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**Search for new heavy narrow resonances
decaying into a dielectron pair with the
CMS detector**

Thèse présentée par

Laurent THOMAS

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Pr. Michel TYTGAT (ULB) (*Président*)

Pr. Barbara CLERBAUX (ULB) (*Promoteur*)

Pr. Gilles DE LENTDECKER (ULB) (*Secrétaire*)

Pr. Freya BLEKMAN (VUB)

Pr. Giacomo BRUNO (UCL)

Pr. Fabienne LEDROIT-GUILLON (Grenoble)

Pr. Claire SHEPHERD-THEMISTOCLEOUS (RAL)

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Résumé

Le sujet de la présente thèse porte sur la recherche de nouvelles particules très massives se désintégrant en une paire électron-positron avec le détecteur CMS.

Le démarrage en 2010 du Large Hadron Collider au CERN marque le début d'une nouvelle ère en physique des particules. L'énergie et l'intensité de ses faisceaux de protons, inégalées à ce jour, offre en effet la possibilité d'étudier les lois décrivant les constituants ultimes de la matière et leurs interactions à des énergies jusqu'alors inaccessibles et d'étudier des processus rares.

La découverte récente par les expériences ATLAS et CMS du boson scalaire prédit par la théorie de la brisure de symétrie électro-faible constitue ainsi la première percée du programme de recherche du LHC et confirme la théorie actuelle décrivant la physique subatomique, le Modèle Standard. Il est cependant largement admis que cette théorie, bien que hautement prédictive et jamais mise en défaut expérimentalement jusqu'à présent, ne constitue qu'une approximation à basse énergie d'une théorie plus fondamentale.

Cette thèse décrit la recherche de nouvelles particules, prédites par plusieurs modèles au delà du Modèle Standard, via leur désintégration en une paire électron-positron de haute énergie. La reconstruction et la sélection des électrons de haute énergie par le détecteur CMS sont des éléments centraux de cette analyse et sont étudiées en détail. Divers critères sont développés afin de distinguer les électrons des autres types d'objets physiques produits lors de collisions de protons, tels que les jets. L'intensité des faisceaux du LHC est telle que plusieurs collisions ont lieu simultanément dans le détecteur et il est montré que l'efficacité de sélection des électrons dépend fortement du nombre de ces interactions. Une technique est donc mise au point pour corriger cet effet.

Une méthode pour mesurer l'efficacité de la sélection directement sur les données est également développée. Celle-ci permet de confirmer les mesures obtenues à partir de simulations, jusqu'à des impulsions transverses de plusieurs centaines de GeV. Le spectre de masse des paires diélectron est établi pour les données enregistrées en 2012 à une énergie dans le centre de masse des protons de 8 TeV, et un excès localisé d'événements est recherché. Aucune déviation significative par rapport au bruit de fonds attendu n'est observée et des limites très contraignantes sont établies sur le rapport de la section efficace de production d'une nouvelle résonance diélectronique et de celle mesurée au pic du boson Z. Ces résultats sont utilisés pour fixer des limites inférieures sur la masse de nouvelles particules prédites par certains modèles.

Le redémarrage du LHC en 2015 avec une énergie de 6.5 TeV par faisceau de proton élargira fortement le potentiel de découverte de ces résonances. En cas de découverte d'un signal, ses propriétés (telles que le spin ou l'asymétrie avant-arrière) seront étudiées avec attention. Des projections sur la précision qui pourrait alors être atteinte pour ces mesures sont donc finalement présentées en fonction de la luminosité intégrée collectée.

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Introduction

“What are we made of ?”. This question has probably raised in everyone’s mind at least once. It is therefore not surprising that it has fascinated thinkers for ages. In the fifth century BC, the Greek philosopher Democritus was amongst the first to propose that Nature could be fully described in term of its ultimate constituents, the smallest indivisible pieces of matter, that he called the atoms. It took more than 2000 years and many efforts during the early twentieth century to acquire the technology to experimentally test this hypothesis and make the discovery of what is now known as the modern atom. However, the indivisibility of the atom did not hold for very long and the twentieth century was a succession of discoveries of new subatomic particles. Although the modern atom is now well known to be divisible, the idea of Democritus of describing the Nature from its smallest constituents has remained.

The current understanding of the ultimate structure of the matter and the way it interacts dates back to the middle of the 1970s and is summarized in a theory called the Standard Model of Particle Physics. This highly predictive theory has never been contradicted by the experience so far, despite many attempts.

The start-up of the Large Hadron Collider in 2009 marks the beginning of a new era in particle physics. Thanks to its intense high energy proton beams it allows one to probe physics in an energy domain out of reach so far. A first milestone was already reached very recently, in 2012, when the last unproven assumption of the Standar Model, the existence of a scalar boson through the interaction of which the elementary particles can acquire mass, was finally confirmed by the ATLAS and CMS experiments at CERN [1, 2].

Yet, despite its impressive successes, there exist strong reasons to believe that the Standard Model is only a good approximation of a more complete theory that could emerge when probing the TeV energy region.

In this thesis, a search for physics beyond the Standard Model is conducted using high energy electron pairs¹ produced in proton-proton collisions at a center of mass energy, $\sqrt{s}$, of 8 TeV and recorded with the CMS detector. The dielectron mass spectrum is scrutinized in the hope of observing a new peak at high mass that would be the sign of new physics. Such a search is motivated by several arguments. First, it takes fully advantage of the unequaled LHC proton beam energy and allows one to probe energies up to several TeV. Moreover, in a hadron collider, the dielectron final state constitutes a clean channel with low background and the signal signature (a new resonance) is very different from what is expected for the background (a sharply decreasing tail). The consequence is that a signal

¹ Unless explicitly stated otherwise, the term “electron” refers to both electron and positron in this thesis.

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can quickly emerge in the data and, in such a case, its existence will rapidly become undoubtful. Finally, there exists a large variety of theoretical models predicting such a signature.

One of the key challenges in this search is the reconstruction and selection of electrons with the CMS detector which must be as efficient as possible on a very large energy range. The development of this selection is one of the focuses of this thesis. A high selection efficiency for very high energy electrons is crucial and a technique to measure it directly on the data was developed.

In the analysis of the 2012 data, no significant deviation to the Standard Model is observed and upper limits are set on the production cross section of new dielectron resonances.

After a first run of 3 years, the LHC is currently under maintenance. This first data taking is only the starting point of a vast research program aimed to last for at least a decade and a project of a major upgrade in 2023 of the LHC and its detectors is already under study, in order to increase the instantaneous luminosity by a factor of 5. The next LHC run, scheduled for 2015 with a center of mass energy raised to 13 TeV will significantly open its discovery potential. It is therefore possible that a signal not detected so far shows up quickly after the LHC resumes. In this eventuality, the improvement offered by the LHC upgrade on the precision that could be reached for the measurement of the properties of a new signal is studied.

This thesis is structured as follows. Chapter 1 presents the current theoretical framework of particle physics, some of its open questions and scenarios of new physics predicting the existence of heavy particles decaying into an electron-positron pair. The Large Hadron collider and the CMS detector that was used to collect the data analyzed in this work are described in chapter 2, as well as the reconstruction of the various types of particles produced during proton-proton interactions. The description of proton-proton collisions and their generation using Monte Carlo simulations is the topic of chapter 3. The identification of high energy electrons in the CMS detector represents an important part of this work and is the focus of the next two chapters: chapter 4 details the various criteria used to select such electrons whereas chapter 5 presents measurements of the efficiency of this selection using data. The results of the search for new heavy resonances decaying into a dielectron pair using the data recorded in 2012 by the CMS detector are then detailed in chapter 6. Finally, chapter 7 presents projections for the search of such resonances at higher energy ($\sqrt{s} = 13$ or 14 TeV) and luminosity, and focuses on the measurement of its properties in case a signal is to be found after the Large Hadron Collider resumes in 2015.

The results presented in chapter 6 are reported as a preliminary public result [3] and a publication is in preparation [4]. The author was also involved in the analysis of the data recorded in 2010 [5] and 2011 [6, 7] at $\sqrt{s} = 7$ TeV. Some of the efficiency measurements presented in chapter 5 have been made public by the CMS Collaboration [8] and will be part of a general publication treating the performances of the electron reconstruction and selection with the CMS detector [9]. The work presented in chapter 7 is expected to be included in a publication detailing the improvement of the physics reach offered by the LHC and CMS upgrades.

Chapter 1

The Standard Model of particle physics and beyond

In this chapter, the theory of the Standard Model of particle physics is first briefly reviewed. Some open questions that motivate the existence of new physics are then detailed. In this thesis, a search for new resonances in the dielectron mass spectrum with the CMS detector is performed. The properties of the Drell-Yan process, which is the main background from the Standard Model in this channel, is presented and various models of new physics leading to high energy dilepton pair production are then described.

1.1 The Standard Model of particle physics

It is currently believed that all the phenomena in Nature can be explained in term of four fundamental interactions between the constituents of matter. The Standard Model [10] is the theory summarizing our current knowledge of the elementary particles of matter and three of their interactions: the electromagnetic and weak interactions, unified to form the electroweak interaction [11, 12, 13], and the strong interaction. The fourth fundamental interaction, gravity, is however left aside.

Mathematically, the Standard Model is a quantum field theory where every matter constituent or interaction is described by one or several fields, each of those being associated to a particle. One therefore distinguishes two kinds of particles: fermions, particles of matter carrying a semi integer spin, and bosons, with integer spins, associated to the interactions.

1.1.1 The fundamental particles of matter

The particles of matter are listed in table 1.1. They carry a spin $1/2$ and are classified into two categories: leptons, only sensitive to the electroweak interaction, and quarks, also sensitive to the strong interaction. Each of these categories is further subdivided into three generations, containing an “up” and a “down” particle. This last feature is connected to the electroweak interaction, as explained below. Up or down particles from different generations only differ by their mass. Finally, each quark exists in three different

1. THE STANDARD MODEL OF PARTICLE PHYSICS AND BEYOND

Leptons			
Neutrinos	ν_e m < 2 eV/c ²	ν_μ m < 2 eV/c ²	ν_τ m < 2 eV/c ²
Charged leptons	e m = 511 keV/c ²	μ m = 106 MeV	τ m = 1.78 GeV
Quarks			
Up quarks	u m = 2.3 MeV	c m = 1.28 GeV	t m = 173 GeV
Down quarks	d m = 4.8 MeV	s m = 95 MeV	b m = 4.7 GeV

Table 1.1: The Standard Model elementary fermionic particles and their measured masses (given here for indication, the exact values and errors can be found in [14]).

colors, a property related to the strong interaction. To each of these particles is associated an anti-particle with same mass and opposite quantum numbers.

1.1.2 Interactions and the Standard Model gauge group

In quantum field theories, the equations of motion of the different fields considered are derived from the action that contains all the information on the field interaction and propagation. For a spin 1/2 particle with mass m , which is described by a spinor field ψ , the action is given in the absence of interaction by:¹

$$S = \int dx^4 \mathcal{L} = \int dx^4 \{ i\bar{\psi}(x)\gamma^\mu \partial_\mu \psi(x) - m\bar{\psi}(x)\psi(x) \} \quad (1.1)$$

where $\mathcal{L}$ is the lagrangian density. The equation of motion leads to the Dirac equation:

$$i\gamma^\mu \partial_\mu \psi(x) = m\psi(x) \quad (1.2)$$

In the Standard Model, the fermionic fields are added by hand to the action to account for experimental observations. The situation is however different for the bosonic fields. The existence of these fields are direct consequences of invariance properties of the action. Symmetries play a central role in physics. In 1918, Emmy Noether indeed stated and proved that to any symmetry of the action is associated a conserved law [15]. The energy

¹ In this thesis, the usual 4-vector formalism is used to describe space and time: the time coordinate and the three spatial coordinates are grouped in a vector $x = x^\mu e_\mu = (x^0, x^1, x^2, x^3)$, where $x^0 = ct$. Repeated indices are summed using the metric tensor $\eta_{\mu\nu}$ that defines the geometry of the Minkowski space-time: $a^\mu b_\mu = \sum_{\mu,\nu} \eta_{\mu\nu} a^\mu b^\nu$, where $\eta_{\mu\nu}$ is defined by $diag(\eta_{\mu\nu}) = (1, -1, -1, -1)$. The following convention is also adopted: 4 vectors are denoted by common type and Greek indices: $a = a^\mu e_\mu$ while 3 vectors are denoted by arrowed type and Latin indices $\vec{a} = a^i \vec{e}_i$.

∂_μ denotes the covariant derivative $\partial_\mu = \frac{\partial}{\partial x^\mu}$ and the γ^μ are a set of four 4-dimensional matrices satisfying the Clifford algebra $\gamma^\mu \gamma^\nu + \gamma^\nu \gamma^\mu = 2\eta^{\mu\nu} I_{4 \times 4}$.

The system of units is chosen such that $\hbar = c = 1$. In this system, space and time quantities have the same dimensions, inverse to energy and momentum quantities.

and momentum conservations are for instance the consequences of the invariance of the action under time and space translations. Another example can be found in the Maxwell equations which unify the electric ($\vec{E}$) and magnetic ($\vec{B}$) fields:

$$\vec{\nabla} \cdot \vec{E} = \frac{\rho}{\epsilon_0} \quad (1.3a)$$

$$\vec{\nabla} \cdot \vec{B} = 0 \quad (1.3b)$$

$$\vec{\nabla} \times \vec{E} = -\frac{\partial \vec{B}}{\partial t} \quad (1.3c)$$

$$\vec{\nabla} \times \vec{B} = \mu_0 \left(\vec{j} + \frac{\partial \vec{E}}{\partial t} \right) \quad (1.3d)$$

where ϵ_0 (resp. μ_0) are the permittivity (resp. permeability) of free space and ρ (resp. $\vec{j}$) are the electric charge (resp. current) densities. One sees from those equations that $\vec{E}$ and $\vec{B}$ can be expressed in term of a single 4 vector $(\phi, \vec{A})$:

$$\begin{cases} \vec{E} = -\vec{\nabla}\phi + \frac{\partial \vec{A}}{\partial t} \\ \vec{B} = \vec{\nabla} \times \vec{A} \end{cases} \quad (1.4)$$

It can be noticed that equations 1.4 remain unchanged under the following transformation:

$$A^\mu(x) \rightarrow A^\mu(x) + \partial^\mu f(x) \quad (1.5)$$

where $f(x)$ can be any function. Such a transformation is called a gauge transformation and can appear naturally in quantum field theory.

Starting from the free spinor action (1.1), one sees that it is invariant under the transformation: $\psi(x) \rightarrow e^{-i\alpha}\psi(x)$. One can maintain this invariance if $\alpha = \alpha(x)$ by adding a new field, $A^\mu(x)$, to the action that replaces the derivative by the covariant derivative:

$$\mathcal{D}_\mu \equiv \partial_\mu + ieA_\mu(x) \quad (1.6)$$

The new action is then given by:

$$S = \int dx^4 \left\{ i\bar{\psi}(x)\gamma^\mu \partial_\mu \psi(x) - e\bar{\psi}(x)\gamma^\mu A_\mu(x)\psi(x) - m\bar{\psi}(x)\psi(x) - \frac{1}{4} (F_{\mu\nu}(x)F^{\mu\nu}(x)) \right\} \quad (1.7)$$

where the last term is the kinematic term of the $A_\mu(x)$ field in which:

$$F_{\mu\nu}(x) \equiv \frac{1}{ie} [\mathcal{D}_\mu, \mathcal{D}_\nu] = \partial_\mu A_\nu(x) - \partial_\nu A_\mu(x) \quad (1.8)$$

The action 1.7 is then invariant under the following set of transformations:

$$\begin{cases} \psi(x) \rightarrow \exp(-ie\alpha(x)) \psi(x) \\ A_\mu(x) \rightarrow A_\mu(x) + \partial_\mu \alpha(x) \end{cases} \quad (1.9)$$

1. THE STANDARD MODEL OF PARTICLE PHYSICS AND BEYOND

Notations	First generation ($i = 1$)	Second generation ($i = 2$)	Third generation ($i = 3$)
Left-handed quarks			
$\Psi_L^{\mathcal{Q}^i}(Y = \frac{1}{6})$	$\begin{pmatrix} (u_{L,r}, u_{L,g}, u_{L,b}) \\ (d_{L,r}, d_{L,g}, d_{L,b}) \end{pmatrix}$	$\begin{pmatrix} (c_{L,r}, c_{L,g}, c_{L,b}) \\ (s_{L,r}, s_{L,g}, s_{L,b}) \end{pmatrix}$	$\begin{pmatrix} (t_{L,r}, t_{L,g}, t_{L,b}) \\ (b_{L,r}, b_{L,g}, b_{L,b}) \end{pmatrix}$
Right-handed quarks			
$\Psi_R^{u^i}(Y = \frac{2}{3})$ $\Psi_R^{d^i}(Y = -\frac{1}{3})$	$\begin{pmatrix} (u_{R,r}, u_{R,g}, u_{R,b}) \\ (d_{R,r}, d_{R,g}, d_{R,b}) \end{pmatrix}$	$\begin{pmatrix} (c_{R,r}, c_{R,g}, c_{R,b}) \\ (s_{R,r}, s_{R,g}, s_{R,b}) \end{pmatrix}$	$\begin{pmatrix} (t_{R,r}, t_{R,g}, t_{R,b}) \\ (b_{R,r}, b_{R,g}, b_{R,b}) \end{pmatrix}$
Left-handed leptons			
$\Psi_L^{\mathcal{C}^i}(Y = -\frac{1}{2})$	$\begin{pmatrix} \nu_{eL} \\ e_L \end{pmatrix}$	$\begin{pmatrix} \nu_{\mu L} \\ \mu_L \end{pmatrix}$	$\begin{pmatrix} \nu_{\tau L} \\ \tau_L \end{pmatrix}$
Right-handed leptons			
$\Psi_R^{l^i}(Y = -1)$	$\Psi_R^e = (e_R)$	$\Psi_R^\mu = (\mu_R)$	$\Psi_R^\tau = (\tau_R)$

Table 1.2: The Standard Model elementary particles of matter classified according to their transformation under the Standard Model gauge group. Each quark type exists in three color states (red, green or blue) and transforms under the triplet representation of $SU(3)_c$. The left-handed particles are assembled in a up ($I_W^3 = 1/2$) and a down ($I_W^3 = -1/2$) particle and transform under the doublet representation of $SU(2)_L$, while right-handed particles ($I_W^3 = 0$) under the singlet representation. Finally, all the multiplets also transform under $U(1)_Y$, with an hypercharge Y indicated in the first column. The fermion electrical charges are obtained from Y and I_W^3 through the relation $Q = Y + I_W^3$. Right-handed neutrinos have never been observed and are therefore not included in this table.

The introduction of a new bosonic field coupling to the matter field is a necessity to achieve the local gauge invariance of the action. In the Standard Model, all the interactions arise in this way: they are consequences of local gauge symmetries.

To construct the invariance gauge group of the Standard Model, one first needs to notice that the 4-dimension spinors can be rewritten as the sum of two semi-spinors:

$$\psi(x) = L\psi(x) + R\psi(x) \equiv \psi_L(x) + \psi_R(x) \quad (1.10)$$

where L (resp. R) is the left-handed (resp. right-handed) projector given by $(1 - \gamma^5)/2$ (resp. $(1 + \gamma^5)/2$). In the Standard Model, the left-handed particles are grouped into doublets consisting of one charged and one neutral lepton or one up and one down quark with a weak isospin, I_W^3 , equal to $\pm \frac{1}{2}$, while right-handed particles are singlets ($I_W^3 = 0$). Each quark in addition forms a triplet of three color states. This structure is illustrated in table 1.2.

The Standard Model action is required to be invariant under $U(1)_Y \times SU(2)_L \times SU(3)_c$, that is, any multiplet $\Psi(x)$ should leave the action invariant under the rotation:

$$\Psi(x) \rightarrow \exp \left(ig \sum_a \alpha^a(x) \frac{\tau^a}{2} \right) \Psi(x) \quad (1.11)$$

where $\alpha^a(x)$ can be any function and the $\frac{\tau^a}{2}$ are the generators of the gauge group representation of $\Psi(x)$: $\tau^a = 0$ for the trivial representations, τ is the weak hypercharge, Y , for $U(1)_Y$, τ^a are the Paulice matrices σ^a ($a=1,2,3$) for $SU(2)_L$ doublets, and the Gell-Mann matrices λ^a ($a=0,\dots,8$) for $SU(3)_c$ triplets. To achieve this, the covariant derivative is therefore defined as¹:

$$\begin{cases} \mathcal{D}_\mu \Psi_L^{\mathcal{Q}^i} \equiv \left\{ \partial_\mu - i\frac{g_1}{2} Y_{\mathcal{Q}^i} B_\mu - i\frac{g_2}{2} \sum_{a=1,2,3} \sigma^a W_\mu^a - i\frac{g_3}{2} \sum_{b=1,\dots,8} \lambda^b G_\mu^b \right\} \Psi_L^{\mathcal{Q}^i} \\ \mathcal{D}_\mu \Psi_R^{u^i, d^i} \equiv \left\{ \partial_\mu - i\frac{g_1}{2} Y_{q^i} B_\mu - i\frac{g_3}{2} \sum_{b=1,\dots,8} \lambda^b G_\mu^b \right\} \Psi_R^{u^i, d^i} \\ \mathcal{D}_\mu \Psi_L^{\mathcal{L}^i} \equiv \left\{ \partial_\mu - i\frac{g_1}{2} Y_{\mathcal{L}^i} B_\mu - i\frac{g_2}{2} \sum_{a=1,2,3} \sigma^a W_\mu^a \right\} \Psi_L^{\mathcal{L}^i} \\ \mathcal{D}_\mu \Psi_R^{l^i} \equiv \left\{ \partial_\mu - i\frac{g_1}{2} Y_{l^i} B_\mu \right\} \Psi_R^{l^i} \end{cases} \quad (1.12)$$

g_1 , g_2 and g_3 being respectively the coupling constants associated to $U(1)_Y$, $SU(2)_L$ and $SU(3)_c$. $\mathcal{Q}^i$ and $\mathcal{L}^i$ label the left-handed quarks and leptons multiplets, u^i and d^i the up quarks and down quarks right-handed multiplets and l^i the lepton right-handed multiplets. The corresponding transformation laws for the gauge fields are given by:

$$\begin{cases} B_\mu \rightarrow B_\mu + \partial_\mu \alpha \\ W_\mu^a \rightarrow W_\mu^a + \partial_\mu \alpha^a + g_2 \sum_{b,c} \epsilon^{abc} W_\mu^b \alpha^c \\ G_\mu^a \rightarrow G_\mu^a + \partial_\mu \alpha^a + g_3 \sum_{b,c} f^{abc} G_\mu^b \alpha^c \end{cases} \quad (1.13)$$

where ϵ^{abc} and f^{abc} are the structure constants of $SU(2)$ and $SU(3)$. The Standard Model action for massless fermions is therefore:

$$\begin{aligned} S = \int d^4x \sum_{\mathcal{Q}^i} i \bar{\Psi}_L^{\mathcal{Q}^i} \gamma^\mu \mathcal{D}_\mu \Psi_L^{\mathcal{Q}^i} + \sum_{u^i} i \bar{\Psi}_R^{u^i} \gamma^\mu \mathcal{D}_\mu \Psi_R^{u^i} + \sum_{d^i} i \bar{\Psi}_R^{d^i} \gamma^\mu \mathcal{D}_\mu \Psi_R^{d^i} \\ + \sum_{\mathcal{L}^i} i \bar{\Psi}_L^{\mathcal{L}^i} \gamma^\mu \mathcal{D}_\mu \Psi_L^{\mathcal{L}^i} + \sum_{l^i} i \bar{\Psi}_R^{l^i} \gamma^\mu \mathcal{D}_\mu \Psi_R^{l^i} \\ - \frac{1}{4} \left(F_{B\mu\nu} F_B^{\mu\nu} + \sum_{a=1,2,3} F_{W\mu\nu}^a F_W^{\mu\nu a} + \sum_{b=1,\dots,8} F_{G\mu\nu}^b F_G^{\mu\nu b} \right) \end{aligned} \quad (1.14)$$

The third line contains the kinematics of the bosonic fields:

$$\begin{cases} F_{B\mu\nu} \equiv \partial_\mu B_\nu - \partial_\nu B_\mu \\ F_{W\mu\nu}^a \equiv \partial_\mu W_\nu^a - \partial_\nu W_\mu^a + g_2 \sum_{b,c} \epsilon^{abc} W_\mu^b W_\nu^c \\ F_{G\mu\nu}^a \equiv \partial_\mu G_\nu^a - \partial_\nu G_\mu^a + g_3 \sum_{b,c} f^{abc} G_\mu^b G_\nu^c \end{cases} \quad (1.15)$$

1.1.3 Spontaneous symmetry breaking

Equation 1.14 does not include any mass term. Because the left-handed fermions transform as doublets and right-handed fermions as singlets under the Standard Model gauge group, a term such as $m \bar{\Psi}_R \Psi_L$ would break the invariance of the equation. This difficulty

¹ From now on, the dependency of the fields and of α in x is made implicit.

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can be eluded by the spontaneous symmetry breaking mechanism proposed by Brout, Englert and Higgs [16, 17, 18] in the following way: one supposes the existence of a new complex scalar doublet field, Φ , that interacts with the fermions and self interacts through quadratic and quartic interactions. This field has $Y_\Phi = \frac{1}{2}$ and is a singlet under $SU(3)_c$. The previous action is therefore completed with the following terms:

$$S_{scalar} = \int dx^4 \sum_{Q^i, w^j} \kappa_{ij} \bar{\Psi}_L^Q \tilde{\Phi} \Psi_R^{w^j} + \sum_{Q^i, d^j} \kappa'_{ij} \bar{\Psi}_L^Q \Phi \Psi_R^{d^j} + \sum_{L^i, l^j} \kappa''_{ij} \bar{\Psi}_L^L \Phi \Psi_R^{l^j} + h.c. \\ + \mathcal{D}^\mu \Phi^\dagger \mathcal{D}_\mu \Phi + \frac{v^2 \lambda}{2} \Phi^\dagger \Phi - \frac{\lambda}{2} (\Phi^\dagger \Phi)^2 \quad (1.16)$$

where $h.c.$ stands for hermitian conjugate, $\kappa_{ij}, \kappa'_{ij}, \kappa''_{ij}$ are complex numbers and $\tilde{\Phi} = i\sigma_2 \Phi^*$. One sees that all these terms are invariant under the Standard Model gauge group. If $v^2 > 0$, the minimum of the potential is reached for $\Phi^\dagger \Phi = v/\sqrt{2}$, that is: the vacuum expectation value of the field is not 0 but $v/\sqrt{2}$. Acquiring this value leads the field to privilege one direction with respect to the three others and spontaneously breaks the symmetry of the action. By convention, the vacuum expected value (v.e.v.) direction is chosen to be:

$$\text{v.e.v.} = \begin{pmatrix} 0 \\ \frac{v}{\sqrt{2}} \end{pmatrix} \quad (1.17)$$

The field then consists of quantum oscillations around this value:

$$\Phi = \begin{pmatrix} \phi^1 + i\phi^2 \\ \frac{v}{\sqrt{2}} + \phi^3 + i\phi^4 \end{pmatrix} \quad (1.18)$$

where ϕ^i ($i = 1, \dots, 4$) are four real scalar fields. The symmetry breaking has several consequences.

- The $\kappa_{ij} \bar{\Psi}_L^i \Phi \Psi_R^j$ terms in the action give rise to mass terms of the type $m_{ij} \bar{\psi}_L^i \psi_R^j$. In the Standard Model, κ'' is diagonal but this is not the case for κ and κ' . The matrices can however be diagonalized by performing a rotation in the generations space. Naming $\tilde{\kappa}, \tilde{\kappa}'$ the corresponding matrices, the up (resp. down) quarks mass eigenstates have then a mass given by $m_{u^i} = \tilde{\kappa}_{ii} \frac{v}{\sqrt{2}}$ ($m_{d^i} = \tilde{\kappa}'_{ii} \frac{v}{\sqrt{2}}$).
- Three of the $SU(2)_L \times U(1)_Y$ gauge bosons acquire also masses. Unlike the fermions, for which the couplings with the scalar field are set by hand in the action, the scalar field couplings to the gauge bosons are directly predicted by the theory. The covariant derivative for the scalar field is:

$$\mathcal{D}_\mu \Phi \equiv \left\{ \partial_\mu - i \frac{g_1}{2} Y_\Phi B_\mu - i \frac{g_2}{2} \sum_{a=1,2,3} \sigma^a W_\mu^a \right\} \Phi \quad (1.19)$$

The $\mathcal{D}^\mu \Phi^\dagger \mathcal{D}_\mu \Phi$ term in the action generates mass terms for the gauge bosons. As in the quark case, a rotation is needed to obtain the mass eigenstates:

$$\begin{pmatrix} Z^\mu \\ A^\mu \end{pmatrix} = \begin{pmatrix} \cos \theta_W & \sin \theta_W \\ -\sin \theta_W & \cos \theta_W \end{pmatrix} \begin{pmatrix} W^{3\mu} \\ B^\mu \end{pmatrix} \quad (1.20a)$$

	Name	Interaction	Spin	Mass
A^μ	Photon	Electromagnetic force	1	$< 10^{-18} \text{ eV}/c^2$
$W^{\pm\mu}$	W bosons	Weak force (charged currents)	1	80.4 GeV
Z^μ	Z boson	Weak force (neutral current)	1	91.2 GeV
$G^{\mu a}_{(a=1,\dots,8)}$	Gluons	Strong force	1	0 (theory)
h	Scalar boson	Mass generation	0	125.5 GeV

Table 1.3: The Standard Model bosonic particles and their mass and spin. The values and limits on the masses are taken from [14]. The quoted mass for the gluon is the theoretical value.

$$\begin{pmatrix} W^{+\mu} \\ W^{-\mu} \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} +1 & -i \\ +1 & +i \end{pmatrix} \begin{pmatrix} W^{1\mu} \\ W^{2\mu} \end{pmatrix} \quad (1.20b)$$

where θ_W is the Weinberg angle and satisfies $\tan \theta_W = g_2/g_1$. In this basis A^μ represents the photon which remains massless, Z^μ the Z boson ($M_Z = \sqrt{g_1^2 + g_2^2}v/2$), and $W^{\pm\mu}$ the W bosons ($M_W = g_2v/2$).

- The fields ϕ^1, ϕ^2, ϕ^4 are absorbed by the three massive gauge fields and provide the degree of freedom corresponding to the longitudinal polarization of the gauge bosons.
- The remaining field, ϕ^3 , acquires a mass and corresponds to a new physical neutral particle denoted h ($Q_h = Y_h + I_{Wh}^3 = 0$).
- The “broken action” is not invariant under $SU(2)_L \times U(1)$ anymore but instead under $U(1)_Q$.

Table 1.3 summarizes the final bosonic content of the Standard Model, after the spontaneous symmetry breaking occurs.

1.1.4 Radiative corrections and renormalisation

The Standard Model action contains all the information needed to compute physical quantities such as decay rates or cross sections. In quantum field theory, the probability of a state $|a\rangle$ to evolve after some time to a state $|b\rangle$ is proportional to the square of the amplitude $\langle b | \hat{S} | a \rangle$, where $\hat{S}$ is called the S-matrix which consists of a time ordered exponential of the interacting hamiltonian derived from the action. It is usually treated perturbatively, that is, the exponential is decomposed into a sequence of terms of increasing powers of the coupling constants. Each of its terms can be described by one or several so-called Feynman diagrams from which an amplitude can be calculated using a set of rules. The leading order represents the classical amplitude and the higher orders are quantum corrections. A common issue when calculating the quantum corrections is the appearance of divergencies. In some theories though, these divergencies can be reabsorbed in the definition of the coupling constants at a given scale through a procedure named renormalisation [19].

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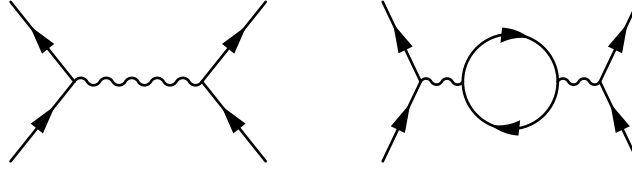


Figure 1.1: The leading order (left) and 1-loop correction (right) Feynman diagrams for the process $\psi\bar{\psi} \rightarrow \gamma \rightarrow \psi\bar{\psi}$ in quantum electrodynamics.

As an example one can consider the action 1.7, where the spinor field is massless. The leading order and the 1-loop Feynman diagrams for the process $\psi\bar{\psi} \rightarrow \gamma \rightarrow \psi\bar{\psi}$ are shown in figure 1.1. The leading order amplitude is straightforward to compute:

$$\mathcal{M}_{tree} = ie^2 \gamma^\mu \frac{\eta_{\mu\nu}}{Q^2} \gamma^\nu \quad (1.21)$$

For the 1-loop contribution, one must integrate over the fermion 4 momentum in the loop, k . The integral is proportional to $\int d^4k/k^4 \sim \int dk/k$. In order to obtain a finite result, one introduces a cutoff Λ inside the integral: $\int^\Lambda dk/k$ and the total amplitude including the 1-loop correction becomes:

$$\mathcal{M}_{1-loop} = ie^2 \gamma^\mu \frac{\eta_{\mu\nu}}{Q^2} \gamma^\nu \left(1 + \frac{e^2}{12\pi^2} \log \frac{Q^2}{\Lambda^2} \right) \quad (1.22)$$

One can get rid of this cutoff by trading the coupling present in the action, e , for the effective coupling e_{eff} at a scale μ . The equation then becomes:

$$\mathcal{M}_{1-loop} = ie_{eff}^2 \gamma^\mu \frac{\eta_{\mu\nu}}{Q^2} \gamma^\nu \left(1 + \frac{e_{eff}^2}{12\pi^2} \log \frac{Q^2}{\mu^2} \right) \quad (1.23)$$

As a consequence, the coupling that should be used to calculate a physical process depends on its scale. The variation of the coupling is described by the renormalisation group equation which in the present case is given, for 1-loop corrections:

$$\frac{d}{d \log Q} e(Q) = \frac{e^3(Q)}{12\pi^2} \quad (1.24)$$

A similar equation holds for any parameter present in the action.

In the Standard Model, the effective couplings g_i ($i = 1, 2, 3$) associated to $U(1)$, $SU(2)$ and $SU(3)$ are constrained by the equations:

$$\frac{d}{d \log Q} g_i(Q) = -\frac{b_i g_i^3(Q)}{(4\pi)^2} \quad (1.25)$$

with [20]:

$$b_1 = -\frac{4}{3}n_g - \frac{1}{10}n_h \quad (1.26a)$$

$$b_2 = \frac{22}{3} - \frac{4}{3}n_g - \frac{1}{6}n_h \quad (1.26b)$$

$$b_3 = 11 - \frac{4}{3}n_g \quad (1.26c)$$

where n_g is the number of generations (3) and n_h the number of scalar bosons (1). One of the important consequences of those equations is the asymptotic freedom of quantum chromodynamics [21, 22]: because of the sign of b_3 , g_3 weakens at high energy and the quarks can then be treated as free particles. This feature plays a crucial role in the calculation of cross sections in hadron collisions.

1.2 A particular Standard Model process: The Drell-Yan process

The leading contribution to the production of high mass dilepton pairs ($M_{ll} > 20$ GeV) comes from a theoretically well known process: the Drell-Yan process [23]. In the Standard Model framework, it is defined as the annihilation of a quark-antiquark pair to a photon or a Z boson decaying into a lepton-antilepton pair and is described at leading order by the two Feynman diagrams drawn in figure 1.2.

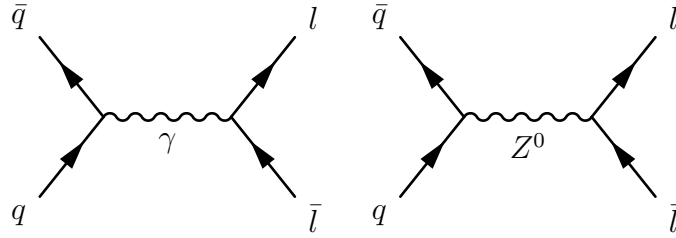


Figure 1.2: Feynman diagrams contributing to the Drell-Yan process at leading order. The left (right) diagram corresponds to the annihilation of a $q\bar{q}$ pair into a photon (a Z boson).

This process is theoretically well known and its signature consisting of two high energy leptons suffers low background in a hadron collider, providing a clean channel to perform precision measurements or to search for hints of new physics.

1.2.1 Partonic cross section

In the $q\bar{q}$ center of mass frame, one defines θ as the angle between the quark and the negative lepton directions of flight. The angular partonic cross section of the Drell-Yan process is then given, at leading order, by¹:

$$\frac{d\sigma_{q\bar{q} \rightarrow \gamma/Z \rightarrow l\bar{l}}}{d\Omega} = \frac{1}{3} \frac{e^4}{(4\pi)^2} Q_q^2 Q_l^2 \frac{1}{4s'} [c_1(1 + \cos^2 \theta) + c_2 \cos \theta] \quad (1.27)$$

¹The detailed calculation is presented in Appendix A.

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where s' is the $q\bar{q}$ invariant mass squared, e is the fundamental charge unit, $Q_{q,l}$ are the quark and lepton electrical charges and c_1 and c_2 are given by:

$$c_1 = 1 + 2\text{Re}(\mathcal{R})g_{V_l}g_{V_q} + |\mathcal{R}|^2(g_{V_l}^2 + g_{A_l}^2)(g_{V_q}^2 + g_{A_q}^2) \quad (1.28a)$$

$$c_2 = 0 + 4\text{Re}(\mathcal{R})g_{A_l}g_{A_q} + 8|\mathcal{R}|^2g_{V_l}g_{A_l}g_{V_q}g_{A_q} \quad (1.28b)$$

with

$$\mathcal{R} = \frac{1}{Q_l Q_q \sin^2 2\theta_W} \frac{s'}{s' - M_Z^2 + i s' \Gamma_Z / M_Z} \quad (1.29a)$$

$$g_{V_{l,q}} = I_{Wl,q}^3 - 2Q_{l,q} \sin^2 \theta_W \quad (1.29b)$$

$$g_{A_{l,q}} = -I_{Wl,q}^3 \quad (1.29c)$$

M_Z and Γ_Z are the Z boson mass and width and $I_{Wl,q}^3$ is the lepton or quark weak isospin. The first, second and third terms in equations 1.28 correspond respectively to the photon, the interference and the Z boson contributions. After integration over $d\Omega$, 1.27 becomes:

$$\sigma_{q\bar{q} \rightarrow \gamma/Z \rightarrow l\bar{l}} = \frac{e^4}{36\pi} Q_q^2 Q_l^2 \frac{1}{s'} c_1 \quad (1.30)$$

$\sigma_{q\bar{q} \rightarrow \gamma/Z \rightarrow l\bar{l}}$ is drawn as a function of $\sqrt{s'}$ in figure 1.3 (left) separately for up and down quarks.

1.2.2 Forward-backward asymmetry

Equation 1.27 is not invariant under $\theta \rightarrow \pi - \theta$. This comes from the fact that right-handed and left-handed particles have different couplings to the Z boson. This feature is quantified with the so-called forward-backward asymmetry A_{FB} defined as¹:

$$A_{FB} \equiv \frac{\sigma_F - \sigma_B}{\sigma_F + \sigma_B} = \frac{3}{8} \frac{c_2}{c_1} \quad (1.31)$$

where $\sigma_F = \sigma_{\theta < \pi/2}$ and $\sigma_B = \sigma_{\theta > \pi/2}$ are the cross sections for which the negative lepton is emitted forward or backward with respect to the quark direction. Figure 1.3 (right) shows the A_{FB} curve for the Drell-Yan process as a function of $\sqrt{s'}$ for up and down quarks. This variable is of special interest in the search for new physics: the angular differential cross section of any spin 1 particle is proportional to $(1 + \cos^2 \theta) + \frac{8}{3} A_{FB} \cos \theta$, A_{FB} can therefore help to distinguish signal from background or to identify a new signal [24]. This will be discussed in section 1.4.2.2 and in chapter 7.

The dilepton mass spectrum predicted by the Drell-Yan process can be divided into different regions depending on which terms dominate in 1.27:

- At low $\sqrt{s'}$ (< 10 GeV) the photon contribution is dominant and the ratio of the $u\bar{u}$ over $d\bar{d}$ cross section reduces to the ratio of their charges squared. The forward-backward asymmetry is zero since the photon couples equally to left-handed and right-handed particles.

¹cf. Appendix A.

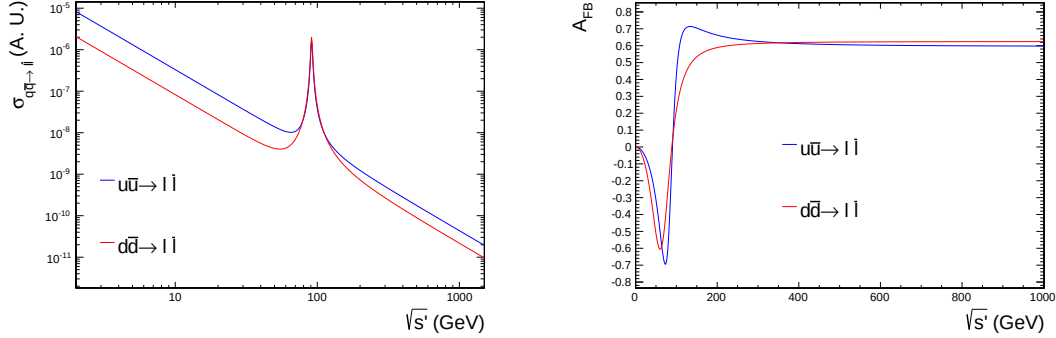


Figure 1.3: Cross section (left) and A_{FB} (right) of the Drell-Yan process at leading order for up (blue) and down (red) quarks.

- For $10 < \sqrt{s'} < 40$ GeV, the effect of the interference, globally negative below the Z mass, is almost not visible in the total cross section but manifests itself in the forward-backward asymmetry.
- For $40 < \sqrt{s'} < 80$ GeV, both the photon, the Z boson and the interference terms are sizeable. The two former ones give the leading contributions to the cross sections whereas the latter is responsible for most of the A_{FB} .
- Between 80 and 100 GeV, the total cross section is completely driven by the Z boson contribution, and the sharp change of A_{FB} with $\sqrt{s'}$ is due to a mix of the Z boson and interference terms. The contribution to the A_{FB} of the Z boson alone is slightly positive and does not depend on $\sqrt{s'}$.
- Above 100 GeV, the three terms have again comparable contributions. As previously, the interference is the main responsible for the non zero A_{FB} whereas the two other terms drive the total cross section.

1.3 Motivations for new physics

Despite the fact that its numerous predictions were experimentally verified with great accuracy in the last 40 years, it is commonly admitted that the Standard Model only constitutes a low energy approximation of a more fundamental theory. Indeed, it does not address several fundamental questions.

- The fourth fundamental interaction, gravity, is not included in the model. Gravity is by many aspects very different from the three other forces and establishing a common framework to describe both faces several difficulties. General relativity indeed indicates that gravity is deeply connected to the space-time geometry and that it couples to the particles energy-momentum tensor, which makes its integration to the Standard Model more subtle than simply adding a new interaction. Moreover,

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whereas the strong, weak and electromagnetic forces have similar strengths at the electroweak scale (energies of ≈ 100 GeV), gravity is more than 10^{30} times weaker and becomes only appreciable for energies around the Planck scale: $E_{Pl} = 10^{19}$ GeV¹. Finally, quantum gravity is known not to be renormalizable: in quantum field theory, the field associated to gravity, the graviton, is a spin 2 particle, $\mathcal{G}^{\mu\nu}$, and one can show that loop corrections including such particles induce divergencies that cannot be reabsorbed through the renormalisation procedure, as opposed to the electroweak and chromodynamics quantum theories.

- A related question is called the fine tuning of the scalar boson mass. Loop corrections to the top quark and W boson mass constraint the scalar boson mass to be $\mathcal{O}(100$ GeV), which is compatible with the newly discovered particle at $M = 125$ GeV. Its tree level mass is however strongly affected by divergent quantum corrections. Top quarks loops lead for instance to a correction given by [25]:

$$\Delta m_S^2 = m_{loops}^2 - m_{tree}^2 = -\frac{|\lambda_t|^2}{8\pi^2} \Lambda_{cutoff}^2 + \dots \quad (1.32)$$

where λ_t is the coupling constant between the top quark and the scalar boson and is close to unity. Λ_{cutoff} is the upper energy bound of the Standard Model domain of validity, E_{Pl} . The correction is therefore proportional to E_{Pl}^2 . In the Standard Model, the complete correction that takes all the particles into account shows the same property. This implies that, for the physical mass to lie in the 100 GeV range, an almost exact cancellation must occur between two $\mathcal{O}(10^{19})$ GeV terms. Avoiding such a miraculous cancellation can only happen if Λ_{cutoff} is $\mathcal{O}(1 - 10)$ TeV rather than the Planck scale, that is, if new physics shows up at that energy. This issue and the previous one are deeply connected and known as the hierarchy problem.

- The structure of the Standard Model is intriguing. First, as mentioned in section 1.1.3, the theory does not explain the values of the fermion masses. All of them are independent and set by hand in the action. The fact that the ratio between the heavier massive fermion (the top quark) and the lighter (the electron) is $\mathcal{O}(10^5)$ seems strange. Another curiosity is related to the so called chiral anomalies cancellations. It can be shown that classical symmetries of the action can be destroyed at the quantum level through loop corrections. This happens in particular for so called chiral symmetries under which left-handed and right-handed fermions transform independently, such as $SU(2)$. The breaking of this symmetry would imply, amongst other consequences, the non conservation of the number of left-handed and right-handed particles separately. However, the symmetry can be maintained at the quantum level if the different loop corrections cancel out. These cancellations set constraints on the gauge group structure of the theory and on the fermions hypercharges. In the Standard Model, this also requires the number of quark and leptons families to be equal. These constraints are verified in the Standard Model but it looks somehow a happy coincidence. It is therefore wondered whether another theory could explain these cancellations in a more natural way.

¹ The Planck scale is the energy associated to the Planck mass defined as $M_{Pl} = \sqrt{\hbar c/G}$, where G is the gravitational constant.

-
- The Standard Model does not explain neutrinos oscillations. Several experiments have indeed observed that neutrinos produced with a given flavour (ν_e , ν_μ , ν_τ) oscillate and acquire another flavour after some time [26, 27]. Theoretically, this can only occur if the neutrinos are massive and the associated mass states (ν_1 , ν_2 , ν_3) do not correspond to the flavour states. The three masses must also be not degenerated. Experimentally, only upper limits on the neutrino masses have been set ($m < 2$ eV). The square mass differences were nevertheless measured and found to be non zero: $\Delta m_{12}^2 = 7.50 \pm 0.20 \times 10^{-5}$ eV², $\Delta m_{23}^2 = 2.32^{+0.12}_{-0.08} \times 10^{-3}$ eV² [14]. Although they are massless in the Standard Model, nothing forbids to add a neutrino mass term in the action. However their small mass values compared to the other fermions might suggest that they acquire it through another mechanism than the electroweak symmetry breaking. The absence of right-handed neutrinos is also striking. It is therefore considered that neutrinos might be their own antiparticle.
 - The Standard Model only describes 5% of the total content of the Universe. This statement results from several recent cosmological observations. First, the measured orbital velocities of stars around their galaxy center [28, 29] is incompatible with the observed matter density in the space: the stars are moving too fast. In order to reconcile the experimental data with the theory, the existence of another kind of matter that does not interact through electromagnetic or strong interactions, the dark matter, has therefore been postulated. A second major result in cosmology is the discovery that the Universe is in accelerated expansion: it was measured that in average galaxies recede from each other and that this escape rate increases with the distance [30, 31]. Such a feature cannot be achieved with usual particles and the Universe is therefore thought to also contain dark energy, a new form of energy with a negative pressure. To account for the experimental observations, dark matter and dark energy have been estimated to compose respectively 25% and 70% of the total Universe content but their fundamental nature remains nowadays a mystery.

A broad research program that involves many experiments all around the world is currently deployed in order to find answers to some of these questions. In this program, the experiments at the Large Hadron Collider (LHC) at CERN play a major role. In particular, thanks to its very intense high energy beams, the new proton-proton accelerator allows one to probe physics at the TeV scale for the first time. As already mentioned, there are strong motivations for new physics to appear in this region and many searches are conducted to bring it out.

1.4 New heavy particles decaying into a dilepton pair

One of the straightforward ways to observe physics beyond the Standard Model at a hadron collider would be through an alteration of the dilepton mass spectrum at high mass. The study of high energy dilepton pairs production is both theoretically and experimentally motivated: there exist many models of new physics giving rise to high energy lepton pairs in the final state [32], and at a hadron collider, the background is low and

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leptons are easy particles to detect and measure. In this section, some models of new physics giving rise to such a final state are described.

1.4.1 Supersymmetry

Supersymmetry [25] is a new type of symmetry of the action where bosons and fermions are considered as components of a common supermultiplet. In this theory, any fermionic (bosonic) field is associated to a supersymmetric bosonic (fermionic) field: any rotation in the supermultiplet space, that is, any mixing of boson and fermion partners, should leave the action invariant. There exist many supersymmetric models. The most simple one, the Minimal Supersymmetric Standard Model (MSSM) associates a superpartner to each Standard Model particle: a spin 0 field (“sfermion”) for the Standard Model fermions, a spin 1/2 field (“gaugino”) for the gauge bosons. In this model, the electroweak symmetry breaking requires the existence of two scalar doublets that will generate 5 physical scalar fields. Their superpartners are called “higgsinos”.

To provide an explanation to the observed mass of the scalar boson, the supersymmetry particles are expected to have masses of the order of 1 TeV, comparable to the electroweak symmetry breaking scale. While the baryon (B) and lepton (L) numbers defined as $B = \frac{1}{3}(n_q - n_{\bar{q}})$ and $L = (n_l - n_{\bar{l}})$ (where n_q , $n_{\bar{q}}$, n_l and $n_{\bar{l}}$ are the number of quarks, antiquarks, leptons and antileptons respectively) are separately conserved in the Standard Model action, this is not true in the MSSM, where term violating B and L can easily appear in the action. To maintain these experimentally verified conservation laws, one defines a new conserved quantum number, the R-parity P_R , as:

$$P_R = (-1)^{(B-L)+2s} \quad (1.33)$$

where s is the particle spin. All Standard Model particles have $P_R = +1$ while all superpartners have $P_R = -1$.

Supersymmetry is extremely attractive in several aspects. First, it provides an explanation to the problem of the fine tuning of the scalar boson mass described in section 1.3. For energies greater than the superpartner masses, the loop corrections to the scalar mass in 1.32 must also take into account the superpartner corrections. The corrections of the bosonic and fermionic superpartners are of different sign and turn out to exactly cancel with the Standard Model corrections. Fixing the scalar boson mass issue implies that the supersymmetric particles cannot be much heavier than 1 TeV. In addition, supersymmetry provides a dark matter candidate: if R-parity is conserved, the lightest superparticle is indeed stable.

Supersymmetric models where R-parity is conserved such as the MSSM forbid the decay of new particles that into a dilepton pair. However, more complicated scenarios do not conserve R-parity [33] and make possible the decay into two leptons of the “sneutrinos” (spin 0 partners of the neutrinos). An interesting property of such scenarios is that the couplings to each generation of lepton are in general not equal. Moreover, extensions of the Standard Model gauge group as discussed in the next section can be easily incorporated. In fact, supersymmetric theories motivate such extensions.

1.4.2 Extension of the Standard Model gauge group

In various new physics scenarios, the Standard Model group is interpreted as the residuum of a larger gauge group that broke down at low energy. Many of them have the appealing feature of unifying the three couplings g_1 , g_2 and g_3 into a single fundamental quantity. The mechanism is analogous to what happens in the electroweak sector where $SU(2)_L \times U(1)_Y$ breaks down to $U(1)_Q$ when the scalar boson acquires its vacuum expectation value. During the process of the fundamental group symmetry breaking, extra $U(1)$ symmetries can appear, introducing new neutral gauge bosons. If their mass lies in the TeV range and they couple to quarks and leptons, it should be possible to produce them and detect them at the LHC. In many models, such particles have a narrow width and their experimental signature will be a new peak in the dilepton invariant mass distribution.

1.4.2.1 The $SU(5)$ model and beyond

This model was initially proposed by Georgi and Glashow in 1974 [34]. It embeds $SU(3)_c \times SU(2)_L \times U(1)_Y$ and the matter particles representations in the following way:

- For each generation, the three color states of the right-handed down quark and the charged and neutral left-handed leptons are grouped into a multiplet transforming according to the $\mathbf{\bar{5}}$ representation of $SU(5)$.
- For each generation, the three color states of the right-handed and left-handed up quarks and of the left-handed down quarks and the right-handed charged lepton are grouped into a multiplet transforming according to the $\mathbf{10}$ representation of $SU(5)$.
- The gauge sector is composed of 24 bosons: the 12 Standard Model gauge bosons, and 12 exotic gauge bosons.
- The unification scale, that is the energy at which the symmetry breaking $SU(5) \rightarrow SU(3)_c \times SU(2)_L \times U(1)_Y$ occurs, is $\mathcal{O}(10^{16} \text{ GeV})$. At that energy, the three Standard Model couplings are unified into a single coupling.

Shortly after it was proposed, this theory acquired a lot of popularity in particular because of its prediction $\sin^2 \theta_W = 0.375$ (at the unification scale) compatible with the measurements at that time. For such a value, the renormalisation equations of the three couplings predict that they converge at $\mathcal{O}(10^{16} \text{ GeV})$ ¹. However latest measurements have significantly lowered the uncertainties and $\sin^2 \theta_W = 0.375$ is now ruled out and consequently g_1 , g_2 and g_3 do not converge, as illustrated on the left-hand side of figure 1.4. This is not true for supersymmetric cases such as the MSSM that can still accommodate the current measurements (right-hand side of figure 1.4).

There exists however another requirement that both supersymmetric and non supersymmetric $SU(5)$ theories do not fulfill: the stability of the proton. In the $SU(5)$ model, the exotics bosons indeed implies that the proton should ultimately decay into a positron

¹In the context of $SU(5)$ a proper normalisation of the generators requires to redefine g_1 as $g_1 = \sqrt{\frac{5}{3}} g_Y$.

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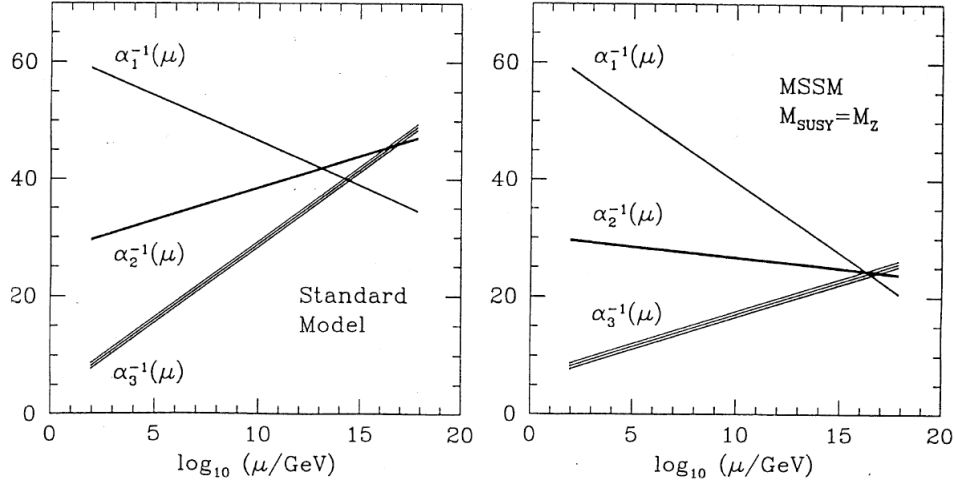


Figure 1.4: Evolution of the couplings $\alpha_1 = \frac{5}{3} \frac{g_1^2}{4\pi}$, $\alpha_2 = \frac{g_2^2}{4\pi}$ and $\alpha_3 = \frac{g_3^2}{4\pi}$ versus the energy scale μ obtained for the renormalisation group equations for the Standard Model (left) and the MSSM (right). While the three couplings do not cross at the same point in the Standard Model, a unification of the couplings is still possible at $\mathcal{O}(10^{16} \text{ GeV})$ in the MSSM. This figure is taken from [35].

and a neutral pion. The reaction is suppressed and the resulting proton half time extremely large ($\approx 10^{29}$ years) but nevertheless smaller than the most recent experimental limits [14]. Finally the $SU(5)$ model cannot include massive neutrinos. In the last decades, the $SU(5)$ model has nevertheless been used as the starting point of a series of new models that could handle the experimental observations. An example is the $SO(10)$ model for which all matter particles belonging to the same generation are grouped into a single multiplet. A nice feature of this model is that the chiral anomalies are here automatically zero for each multiplet, as opposed to $SU(5)$ or the Standard Model in which the cancellations result from the contribution of particles transforming according to different representations.

In the $SU(5)$ model, the new physics scale is $\mathcal{O}(10^{16} \text{ GeV})$, and so is the mass of the new gauge bosons. There is therefore no hope to produce them at the LHC. It can however happen that the symmetries are broken at different scales for larger groups. For instance, the breaking scheme for $SO(10)$ is:

$$SO(10) \rightarrow SU(5) \times U(1)_\chi \rightarrow G_{SM} \times U(1)_\chi \quad (1.34)$$

where $G_{SM} = SU(3)_c \times SU(2)_L \times U(1)_Y$ and χ is the charge associated to the extra $U(1)_\chi$ group. There is no constraint on the breaking scale for this $U(1)_\chi$ and it might happen at the TeV scale.

1.4.2.2 The E_6 models

One of the most popular gauge groups embedding $SU(5)$ is E_6 . It can be broken down to the Standard Model group through the following scheme:

$$E_6 \rightarrow SO(10) \times U(1)_\psi \rightarrow SU(5) \times U(1)_\psi \times U(1)_\chi \rightarrow G_{SM} \times U(1)_\psi \times U(1)_\chi \quad (1.35)$$

Model	c_V^u	c_A^u	c_V^d	c_A^d	c_V^l	c_A^l
Z'_ψ	0	0.300547	0	0.300547	0	0.300547
Z'_η	0	0.380165	-0.285124	0.095041	0.285124	0.095041
Z'_χ	0	0.073458	-0.416249	-0.342792	0.416249	-0.342792
Z'_I	0	0	0.620752	-0.620752	-0.620752	-0.620752
Z'_{SSM}	-0.227388	0.592979	0.410183	-0.592979	0.044592	-0.592979

Table 1.4: Vector (c_V) and axial-vector (c_A) couplings of the Z' boson to up quarks (u), down quarks (d) and the charged leptons (l) for various models corresponding to different values of θ_{E_6} . The values are calculated according to the convention adopted in [24]. For comparison, the Z'_{SSM} couplings, which are identical to the Z boson couplings are also given.

It is generally assumed that only one linear combination of the $U(1)_\psi$ and $U(1)_\chi$ Z' s is relevant at the TeV scale:

$$Z' = Z'_\psi \cos \theta_{E_6} + Z'_\chi \sin \theta_{E_6} \quad (1.36)$$

Each θ value gives rise to a specific model, with its own couplings to the fermions. Some specific models are:

- The Z'_ψ ($\theta_{E_6} = 0$) that only interacts through axial-vector couplings with the fermions.
- The Z'_χ ($\theta_{E_6} = \pi/2$).
- The Z'_η ($\theta_{E_6} = \arctan \sqrt{3/5}$), suggested by superstring theories [36].
- The Z'_I ($\theta_{E_6} = \arctan -\sqrt{5/3}$), that does not couple to up quarks and only couples to left-handed down quarks and right-handed leptons.

The Z' couplings to up quarks, down quarks and charged leptons are given in table 1.4 for those models. A property shared by any Z' from E_6 is the absence of vectorial coupling with the up quarks. The couplings of another popular model, the sequential to Standard Model Z' (Z'_{SSM}), which consist of a heavy copy of the Z boson are also presented. Each set of couplings induces a particular forward-backward asymmetry, as illustrated in figure 1.5 [24], which shows the A_{FB} as a function of the dilepton mass for various Z' from the E_6 model with a mass of 500 GeV. For a Z'_ψ which couples the same way to left-handed and right-handed particles, the asymmetry is null at the resonance mass.

1.4.3 Extra dimensions scenarios

Another sort of theories beyond the Standard Model are those postulating the existence of new compactified spatial dimensions. They attempt to explain the apparent weakness of gravity by assuming that the spin 2 field associated to gravity, the graviton, can propagate into these new dimensions as opposed to the Standard Model matter fields whose wave functions are only sizeable in the usual 4-dimensional space-time. Gravity is perceived

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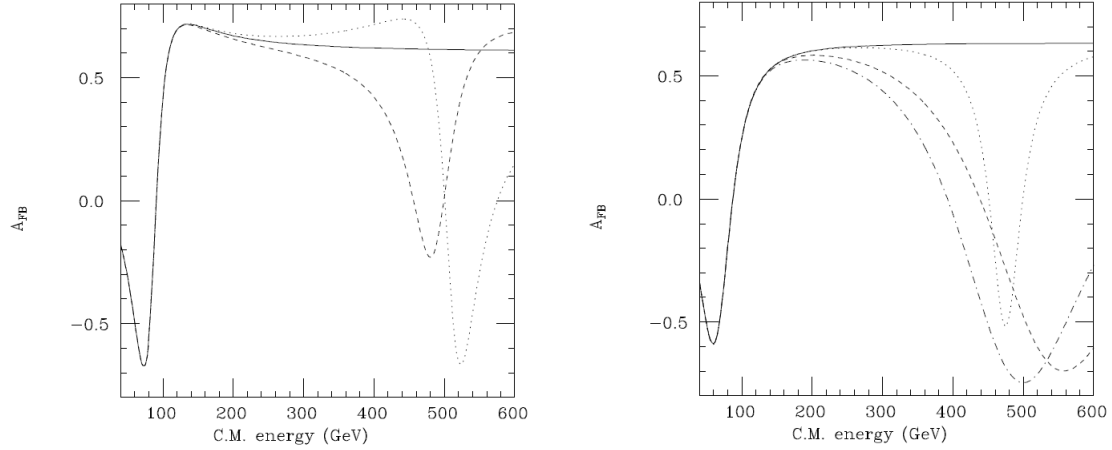


Figure 1.5: The forward-backward asymmetry for the $u\bar{u} \rightarrow l\bar{l}$ (left) and $d\bar{d} \rightarrow l\bar{l}$ (right) processes, as a function of the center of mass energy for the Standard Model (solid line) and various models of Z' at $M=500$ GeV from the E_6 model: Z'_ψ (dotted line), Z'_χ (dashed line) and Z'_I (dot-dashed line). The latter does not couple to up quarks and is therefore absent in the left plot. This figure is taken from [24].

as weak in the 4-dimensional world because of the small overlap between the Standard Model particles and graviton wave functions.

1.4.3.1 Large extra dimensions and ADD gravitons

The fact that the gravitational strength between two masses decreases as the inverse of their distance squared is a consequence of the existence of three spatial dimensions. Whereas this dependency is very well established for macroscopic and cosmological distances it has never been tested for distances below $100 \mu\text{m}$. In 1999, Arkani-Hamed, Dvali and Dimopoulos (ADD) therefore proposed the existence of new extra dimensions with a radius smaller than $100 \mu\text{m}$ [37]. In the case of n extra dimensions with a same radius R , the $1/r$ dependency of the gravitational potential is replaced by a $1/r^n$ dependency for distances $\ll R$. For distances $\gg R$, the $1/r$ law holds but the usual potential given by:

$$V(r) \sim \frac{m_1 m_2}{M_{Pl}^2} \frac{1}{r} \quad (1.37)$$

is replaced by:

$$V(r) \sim \frac{m_1 m_2}{M_{Pl(4+n)}^{n+2}} \frac{1}{R^n r} \quad (1.38)$$

The effective 4-dimensional Planck mass is therefore related with the intrinsic Planck mass in $4+n$ dimensions:

$$M_{Pl}^2 \sim M_{Pl(4+n)}^{n+2} R^n \quad (1.39)$$

The hierarchy problem is thus eluded if $M_{Pl(4+n)} \sim 1 \text{ TeV}$ which turns into a condition on the extra dimensions size:

$$R \sim 10^{\frac{30}{n}-19} \text{ m} \sim 10^{\frac{30}{n}-12} \text{ eV}^{-1} \quad (1.40)$$

n	R (m)	R (eV $^{-1}$)
1	10^{11} m	$5 \cdot 10^{17}$ eV $^{-1}$
2	$100 \mu\text{m}$	0.5 meV $^{-1}$
3	1 nm	5 keV $^{-1}$
4	3 pm	15 MeV $^{-1}$
5	100 fm	0.5 MeV $^{-1}$
6	10 fm	0.05 MeV $^{-1}$

Table 1.5: Radius of the extra dimensions in the ADD model to solve the hierarchy problem as a function of the number of extra dimensions. The radius is expressed in distance and energy units.

Table 1.4.3.1 shows the corresponding extra dimension size for 1 to 6 extra dimensions.

The $n = 1$ case would lead to $R \sim 10^{11}$ m and is excluded but other values are still possible. ADD dimensions are called large extra dimensions because their radius is in the range $1 \text{ meV}^{-1} - 1 \text{ MeV}^{-1}$ for $n \leq 6$, well above the TeV length $l_{\text{TeV}} = 1 \text{ TeV}^{-1}$. Phenomenologically, ADD dimensions might manifest at the LHC through the production of gravitons. Those particles can propagate in the ADD dimensions and escape our 4-dimensional world, inducing a large amount of missing energy in the detector. Virtual gravitons might also be produced. They have similar couplings to all Standard Model particles and one of the easiest channels to detect them is their decay into a dielectron or dimuon pair. The compactification of extra dimensions implies for particles that can travel into them the existence of a serie of excited states called Kaluza-Klein towers. The mass difference Δm between two consecutive states is

$$\Delta m \sim 1/R \quad (1.41)$$

Each state has a coupling of $\mathcal{O}(\text{Energy}/M_{Pl})$ but because the ADD dimensions are large, Δm is tiny for the energy range probed at the LHC and the number of states involved is huge. As a consequence, the experimental signature of ADD gravitons is not a serie of resonances but rather a general excess in the high mass dielectron and dimuon mass spectra due to the effect of a continuum of graviton states.

1.4.3.2 Kaluza Klein excitations of the gauge bosons

In addition to large extra dimensions described in the previous section, it is possible to imagine other smaller extra dimensions, with $R \sim 1 \text{ TeV}^{-1}$, where the graviton but also the Standard Model fields can propagate. The size of the dimension induce a very different phenomenology: here the Kaluza-Klein modes are spaced every few TeV, and only the first one is likely to be visible at the LHC.

1.4.3.3 Warped extra dimensions: the Randall-Sundrum model

Another type of extra dimension scenario has been developped by Randall and Sundrum [38]. They pointed out that while the ADD model solves the hierarchy problem by

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lowering M_{Pl} to the electroweak scale, a new large hierarchy is introduced between the electroweak scale and the invert of the extra dimension radius R . As an alternative, they considered a model with one extra dimension with $1/R = 10^{19}$ GeV. This fifth dimension is labelled by the coordinate $\phi \in [-\pi, \pi]$, whose boundaries conditions are periodic and of period π and the usual 4-dimensional Universe is the 3-brane (the “slice”) of the five dimensional space-time located at $\phi = \pi$. The 5-dimensional action for the gravity sector is given by:

$$S_{gravity} = \int d^4x \int_{-\pi}^{\pi} d\phi \sqrt{-G} \{ -\Lambda + 2M_{Pl,5d}^3 R \} \quad (1.42)$$

where G is the determinant of the 5-dimensional metric $G^{\mu\nu}$, Λ the cosmological constant, $M_{Pl,5d}$ the 5-dimensional Planck mass, and R the 5-dimensional Ricci scalar. It can be shown that if one requires the solution of five-dimensional Einstein’s equations to satisfy the four-dimensional Poincaré invariance, the 5-dimensional metric is:

$$ds^2 = e^{-2kR|\phi|} \eta_{\mu\nu} dx^\mu dx^\nu + R^2 d\phi^2 \quad (1.43)$$

This metric has two important consequences:

- The 5-dimensional and 4-dimensional Planck masses are related by:

$$M_{Pl}^2 = \frac{M_{Pl,5d}^3}{k} (1 - e^{-2kR\pi}) \quad (1.44)$$

where k is an arbitrary scale. They are therefore close to each other for $k \sim R$.

- The electroweak energy scale v or any physical mass m in the $\phi = \pi$ 3-brane is exponentially suppressed compared to the 5-dimensional energy v_0 or mass m_0 :

$$m = e^{-kR\pi} m_0; v = e^{-kR\pi} v_0 \quad (1.45)$$

In the Randall-Sundrum model, the only fundamental scale is the Planck scale: $v_0 \sim k \sim 1/R \sim M_{Pl,5d} \sim E_{Pl} \sim 10^{19}$ GeV. Bringing the electroweak scale down to 1 TeV in our 4-dimensional Universe can be achieved for $kR \approx 50$ and does not introduce any new hierarchy.

The phenomenological implications of this scenario is the existence in our 4-dimensional Universe of graviton resonances with properties characterized by the electroweak scale in our 4-dimensional Universe¹. The mass difference between two consecutive states is therefore expected to be around 1 TeV, and the coupling of each state is $\mathcal{O}(\text{Energy}/1 \text{ TeV})$. The Randall-Sundrum graviton signature is therefore a narrow resonance in the high tail of the dilepton mass spectrum.

1.4.4 Status of the searches before the LHC era

The hunt for new heavy particles that could decay into a dielectron pair did not start at the LHC. There exist various other experiments that have already set constraints on their properties.

¹ The experimental couplings of this graviton to the Standard Model particles are completely determined by the quantity c , defined as $c = k/M_{Pl,5d}$.

Model	Z'_χ	Z'_ψ	Z'_η	Z'_{SSM}
Lower limit on the mass (GeV)	673	481	434	1787

Table 1.6: The 95% confidence level lower limits on the Z' mass for various models obtained by the analysis of the LEP 2 data [42].

- Precise cross section measurements at the Z pole are sensitive to the existence of a new spin 1 particle, Z' , with a significant mass mixing with the Z boson. Results from LEP1 have however not shown any deviation with respect to the Standard Model expectation and were able to constraint the absolute value of a mixing angle with an hypothetical Z' to be below 0.01 radian for many models [39].
- The combination of Z pole precision observables obtained at LEP, SLC or Tevatron with other electroweak observables which include results from neutrino scattering or atomic parity violation experiments allows one to set lower limits on the mass of the Z' typically between a few hundreds of GeV up to 1 TeV [40, 41].
- Indirect constraints can also be derived from the study of electron-positron collisions with a center of mass energy beyond the Z pole but below the Z' mass. At the LEP 2, the leptonic and hadronic cross sections and the forward-backward asymmetry were used to constraint Z' models [42]. The corresponding lower limits on the mass of those Z' are shown in table 1.6.
- Finally, direct searches in the high mass tail of the Drell-Yan process have already been conducted at the Tevatron. The CDF and D0 collaborations excluded a Z'_{SSM} in the dielectron channel below 963 GeV and 1023 GeV respectively [43, 44]. CDF also set a lower limit of 1058 GeV on the mass of a Randall-Sundrum graviton ($c=0.1$) (combination of dielectron and diphoton channel) [45].

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Chapter 2

The experimental setup

Since the Standard Model theory was proposed in the 1970s, a large program has been deployed to test its predictions. Several generations of colliders have been built with beam energies and intensities constantly increasing while the detectors became more and more sophisticated. This program led to many breakthroughs: the UA1 and UA2 experiments at the superproton-antiproton synchrotron (Sp $\bar{p}$ S) discovered the W and Z bosons [46, 47, 48, 49], the Large Electron-Positron collider (LEP) allowed one to confirm the Standard Model predictions with great accuracy, especially in the newly discovered electroweak sector [50, 51] and the CDF and D0 detector at the proton-antiproton accelerator at FermiLab, the Tevatron, confirmed the existence of the top quark [52, 53]. At the beginning of the Large Hadron Collider (LHC) era, only one piece of the puzzle was missing: the experimental evidence of the scalar boson. The machine was specifically designed to rapidly bring evidence of this particle or alternatively to undubitably exclude its existence. Three years after the LHC start-up, the existence of a scalar boson is now well established.

As advertised in the previous chapter, despite all these successes, the Standard Model is nevertheless considered as an effective theory of a more fundamental theory of whose some hints could be found at the TeV scale. Thanks to its huge center of mass energy and instantaneous luminosity, the second goal of the LHC is therefore the search for new physics.

This chapter describes the experimental apparatus used to provide and collect the data studied in this thesis: the Large Hadron Collider is first briefly presented and the CMS detector, that plays a central role in this thesis, is then detailed.

2.1 The Large Hadron Collider

The Large Hadron Collider [54] is the new hadron collider located at CERN in Geneva. It can generate three types of collisions: symmetric proton-proton or lead-lead collisions, or asymmetric proton-lead collisions. The machine consists of a 27 km ring located 100 meters underground in the tunnel that was built for the LEP in 1988. Its construction, decided in 1994, was completed about 15 years later and the first collisions were recorded in 2009 by the four dedicated experiments: ATLAS [55], ALICE [56], CMS [57] and LHCb [58].

2. THE EXPERIMENTAL SETUP

2.1.1 Machine overview

The LHC is designed to accelerate protons or ions up to an energy of 7 TeV or 2.75 TeV per nucleon respectively. As this thesis presents a search for new physics in proton-proton collisions, this chapter focuses on this type of collisions. The motivations for colliding proton beams at the LHC, rather than electrons with positrons or protons with antiprotons are of two different kinds. Electrons lose a lot of energy through synchrotron radiation when they are deflected, and one quickly reaches a point where the radiofrequency cavities are not able to provide more energy than what is dissipated in one revolution. The radiated power is proportional to the fourth power of the accelerated particle energy over mass ratio and the effect is attenuated by a factor larger than 10^{12} for protons for which the experimental limiting factor becomes the intensity of the magnetic field. This explains why, with the same radius for the ring, the maximal energy reached at LEP was 104.5 GeV per electron while LHC is able to accelerate protons up to 7 TeV. Furthermore, protons are much easier to produce than antiprotons. Working with two protons beams allows one to reach higher beam intensities and to significantly increase the number of collisions per second with respect to the proton-antiproton case.

The proton acceleration is achieved in several steps: protons are extracted by ionizing hydrogen atoms and first accelerated in a linear accelerator, the LINAC2. They then enter a succession of intermediate synchrotrons of increasing size (Proton Synchrotron Booster (PSB), Proton Synchrotron (PS), and Super Proton Synchrotron (SPS)) where their energy is gradually increased to reach 450 GeV. The LHC finally brings the beams to the nominal energy and, in order to increase the collision probability between the proton of each beam, the transverse size of the beam is squeezed before the two beams finally collide. A similar chain holds for lead ions. The complete acceleration phase takes between 1 and 2 hours and a usual collision run lasts for around 10 hours after which the beams are dumped.

The accelerator complex is shown in figure 2.1 and a schematic view of the LHC ring is drawn in figure 2.2, together with the position of the four main experiments and the areas where the beams are injected and dumped.

In the LHC, protons or lead ions are accelerated every revolution thanks to 16 radiofrequency cavities that can generate a maximum voltage of 2 MV each. In order to maintain the particles on a circular trajectory, the LHC ring is made of a succession of thousands of magnets that deflect (dipoles) and collimate (quadrupoles and higher order multipoles) the beams. For a relativistic particle of charge q describing a uniform circular motion, the relationship between its momentum p , its curvature radius R and the magnetic field B is:

$$p = qBR \tag{2.1}$$

The magnetic field needed to achieve the targeted energy is therefore a couple of Teslas. Such an intense field is obtained by using magnets cooled down to 1.9K with liquid Helium, that allows them to work in a superconductor regime and produce a magnetic field of 8 T.

The beams themselves are constituted of bunches of 10^{11} protons around 30 cm long and with a transverse size of the order of a fraction of a millimeter. The two beams travel in opposite directions in two parallel pipes, each 6.3 cm in diameter, maintained at a pressure of 10^{-13} bar. At its nominal capacity, the machine contains 2808 bunches and

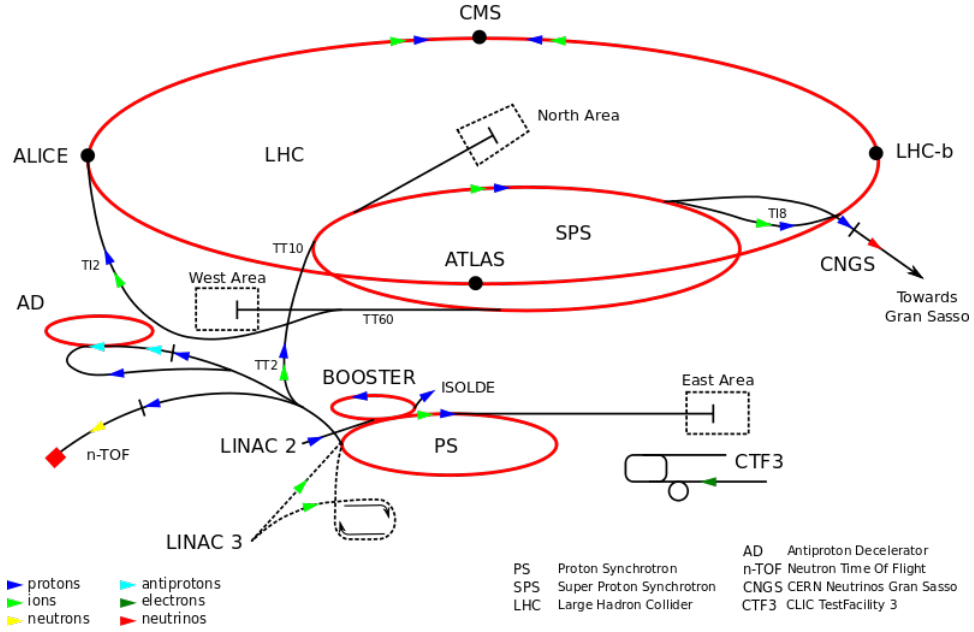


Figure 2.1: The LHC accelerator complex. Before entering the LHC ring, the protons or ions are first accelerated by a series of preaccelerators.

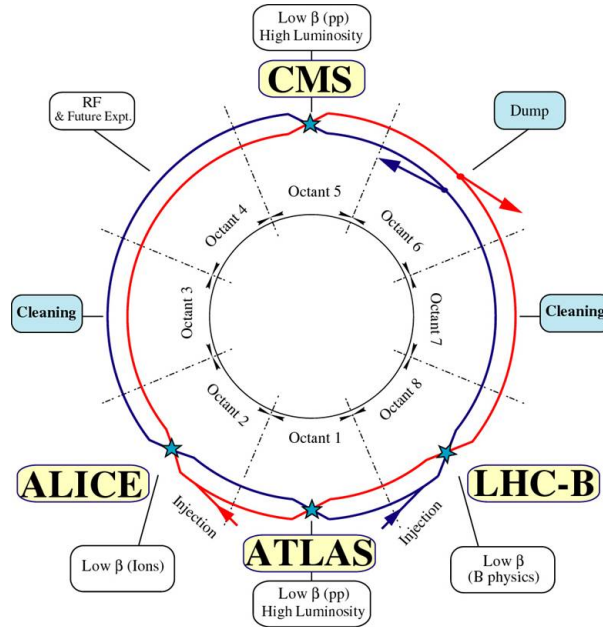


Figure 2.2: Schematic view of the LHC ring. The two beams are collided in 4 interaction points corresponding to the position of the 4 main LHC experiments.

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produce collisions every 25 nanoseconds.

An important characteristic of an accelerator is its instantaneous luminosity, $\mathcal{L}$, which relates the cross section of a given process σ to the experimental event rate $\frac{dN}{dt}$:

$$\frac{dN}{dt} = \sigma \mathcal{L} \quad (2.2)$$

The luminosity depends on the experimental setup. In the case of head-on collisions of bunches of ultrarelativistic particles, it is given by [59]:

$$\mathcal{L} = \frac{n_1 n_2 f}{4\pi \sigma_x \sigma_y} \quad (2.3)$$

where n_1, n_2 are the number of particles per bunch in each beam, f the collision frequency, and σ_x, σ_y the transverse widths in x and y of the beams. The nominal LHC value is $\mathcal{L} = 10^{34} \text{ cm}^{-2}\text{s}^{-1}$, which is a factor 50 higher than the luminosity reached at the Tevatron. At the LHC, the luminosity is directly measured by the different experiments. The CMS detector uses an online measurement, based on the energy deposits in the hadronic forward calorimeters [60]. The results are cross checked with two offline measurements, using the forward calorimeters and tracking information.

2.1.2 The LHC experiments

The LHC research program involves over 10000 scientists from more than 100 different countries. It covers most of the topics of modern particle physics by using information collected by the following experiments:

- ATLAS (A Toroidal LHC ApparatuS) and CMS (Compact Muon Solenoid) are the two biggest experiments. The two detectors were built using different technologies but have similar physical performance. They have a geometric coverage of almost 4π and are designed to handle high instantaneous luminosities. They can therefore fully exploit the LHC potential and look for very rare processes to achieve their main goals: the search for the scalar boson and for physics beyond the Standard Model. Both experiments are multipurpose detectors: in addition to these analyses, they also perform a wide range of other studies in related topics (electroweak and quantum chromodynamics measurements, top quark physics,...).
- The LHCb (Large Hadron Collider beauty experiment) aims to shed some light on CP violation by making precision measurements in the heavy flavour physics sector, and to measure rare decays of B hadrons. It consists of a single-arm spectrometer which, as opposed to CMS and ATLAS, only covers a small solid angle, in the forward region. One of its main features is a high resolution tracking system extended to very forward regions in order to detect B hadrons before they decay. In order to reduce the radiation damage to the subdetectors and the average pile-up which strongly increases the tracker occupancy, the experiment operates at an instantaneous luminosity around ten times lower than ATLAS and CMS.

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- ALICE (A Large Ion Collider Experiment) focuses on the study of the strong interaction and of the quark-gluon plasma, a state of matter that was thought to populate the Universe during the first microseconds after the Big Bang, and that can be reproduced in high energy ion collisions. The detector is made of a barrel part composed of a tracker and calorimeters, and a single arm forward muon spectrometer. It is designed to handle the extreme particle multiplicity of ion collisions.
 - In addition to the four detectors, two smaller experiments are installed: TOTEM (TOTal cross section, Elastic scattering and diffraction dissociation Measurement at the LHC), installed next to the CMS detector, which is a very forward detector dedicated to the measurement of the total proton-proton cross section and the study of elastic and diffractive scattering, and LHCf (Large Hadron Collider Forward) consisting of two forward calorimeters placed on either side of the ATLAS experiment that aims to measure the production of neutral particles at very low angles.

2.1.3 Data taking conditions during Run 1 (2010-2012)

The first LHC collisions were recorded on November 20, 2009 at the injection energy of 450 GeV per beam while the first collisions of protons accelerated by the machine at an energy of 3.5 TeV per beam were recorded on March 30, 2010. This marked the beginning of the LHC Run 1 that extends from 2010 to 2012. During the years 2010 and 2011, the machine operates at a proton-proton center of mass energy of 7 TeV. In 2012, the energy was raised to 8 TeV. As illustrated in figure 2.3, the instantaneous luminosity has continuously increased throughout Run 1 to reach a peak value of $7.7 \times 10^{33} \text{ cm}^{-2}\text{s}^{-1}$ in the last months of 2012. This was achieved despite that unstability issues of the beams forced to reduce the number of bunches to half the nominal value. The total integrated luminosity delivered by the machine to the CMS detector, defined as $\int \mathcal{L} dt$, is shown in figure 2.4. It increased from 44.2 pb^{-1} in 2010 to 6.1 fb^{-1} in 2011 and 23.3 fb^{-1} in 2012. The total luminosity recorded by the detector and used for physics analyses was around 90% of the delivered one, leading to a integrated luminosity of $19.7 \pm 0.5 \text{ fb}^{-1}$ for the 2012 data taking [61]. Because of the high intensity of the LHC beams, several collisions occur in average at each bunch crossing and all of them are recorded at the same time by the detectors. This phenomenon, called pile-up, introduces noise: in addition to an interesting interaction, one must also deal with others that are superimposed. Figure 2.5 presents the distribution of the mean number of interactions per bunch crossing in 2012. That year, an average number of 21 collisions was recorded per bunch crossing. Pile-up is therefore a sizeable effect and its correct treatment represents an important challenge at the LHC.

2.1.4 Expected data taking conditions for the next runs

After a technical stop in 2013 and 2014, the LHC will restart in the first months of 2015 with an energy of 6.5 TeV per beam [62]. The LHC phase 1 is expected to operate until 2022. Beyond 2022, a project to upgrade the machine in order to increase the instantaneous luminosity to five times its nominal value is currently under study. This

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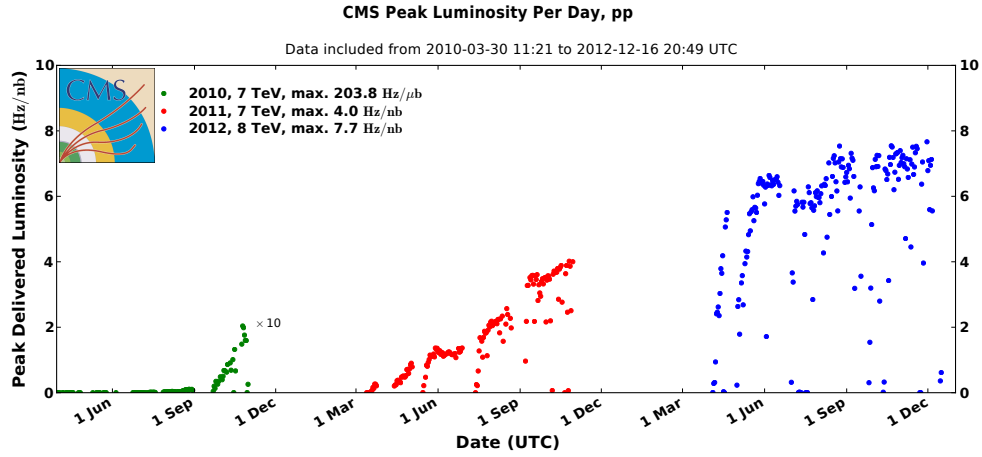


Figure 2.3: Peak instantaneous luminosity per run as a function of time for the three first years of the LHC operations.

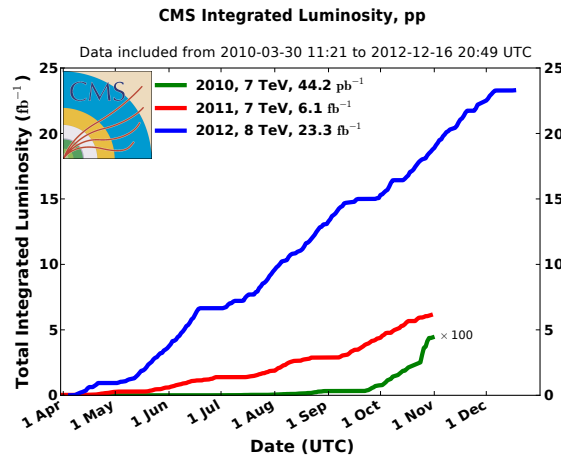


Figure 2.4: Integrated luminosity as a function of time for the three first years of the LHC operations.

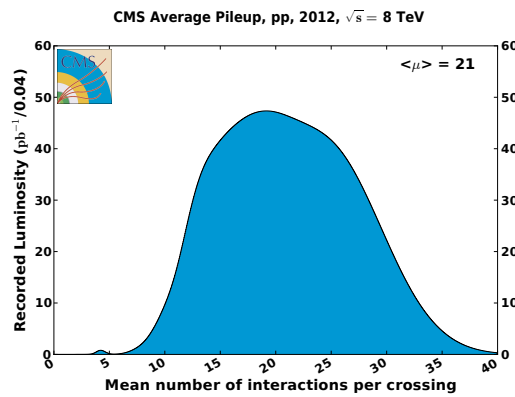


Figure 2.5: Mean number of interactions per bunch crossing recorded by CMS in proton-proton collisions in 2012.

Name	Period	$\sqrt{s}$	Peak luminosity	$\int \mathcal{L} dt$	bunch spacing	Pile-up
Run 1	2010-2012	7-8 TeV	$7.7 \cdot 10^{34} \text{ cm}^{-2}\text{s}^{-1}$	25 fb^{-1}	50 ns	20
Run 2	2015-2017	13-14 TeV	$1.6 \cdot 10^{34} \text{ cm}^{-2}\text{s}^{-1}$	100 fb^{-1}	25 ns	40
Run 3	2019-2021	14 TeV	$2 \cdot 10^{34} \text{ cm}^{-2}\text{s}^{-1}$	300 fb^{-1}	25 ns	50
Phase 2	2023- $\approx$ 2035	14 TeV	$5 \cdot 10^{34} \text{ cm}^{-2}\text{s}^{-1}$	3000 fb^{-1}	25 ns	140

Table 2.1: The expected LHC operations schedule and the data taking conditions from the LHC start-up in 2010 until 2035.

second phase is foreseen to extend the LHC program until 2035. Its full schedule [63] is detailed in table 2.1. During the phase 1, the data taking periods are subdivided into three runs, between which the beam energy and instantaneous luminosity will be increased, as shown in table 2.1 where various other data taking parameters are also given. The run 1 corresponds to the 2010-2012 period during which data corresponding to an integrated luminosity of 25 fb^{-1} were collected. The expected luminosities at the end of run 2 and run 3 are 100 and 300 fb^{-1} respectively. The full luminosity collected by the end of the LHC phase 2 is estimated to reach 3000 fb^{-1} .

During the long shutdown periods separating the three phases, an important upgrade program for the accelerators and the detectors will be deployed. These upgrades are especially needed to reach and sustain the targeted instantaneous luminosity for the LHC phase 2 ($5 \cdot 10^{34} \text{ cm}^{-2}\text{s}^{-1}$). The high radiation level will in addition cause aging of some detector components. By the beginning of phase 2, most of the detector components currently in operation will have been replaced by new ones with improved performance.

2.2 The CMS experiment

The Compact Muon Solenoid (CMS) is one of the two multipurpose detectors at the LHC. It has the form of a 28.7 meter long cylinder, with a diameter of 15.0 meters. Its total weight of 14000 tonnes makes it the heaviest detector ever built at a particle collider. This section presents its substructure and its experimental performance as well as the triggering system.

2.2.1 Overall structure

The detector is made of a succession of layers of subdetectors with different purposes. Starting from the interaction point, the particles first pass through the tracking system where the trajectory of charged particles, which ionize the medium, can be reconstructed. They then enter the calorimeters, which absorb and measure the energy of photons and electrons (electromagnetic calorimeter) and hadrons (hadronic calorimeter). Those components are embedded in a solenoid producing a magnetic field of 3.8 T parallel to the beam axis, providing a way to measure the momentum of charged particles. Finally, muons chambers interleaved with a steel return yoke are placed around the solenoid.

The detector is divided in a cylindric central part, the barrel, closed by two endcaps. Because of their different radiation fluxes and geometry, different technologies are used in

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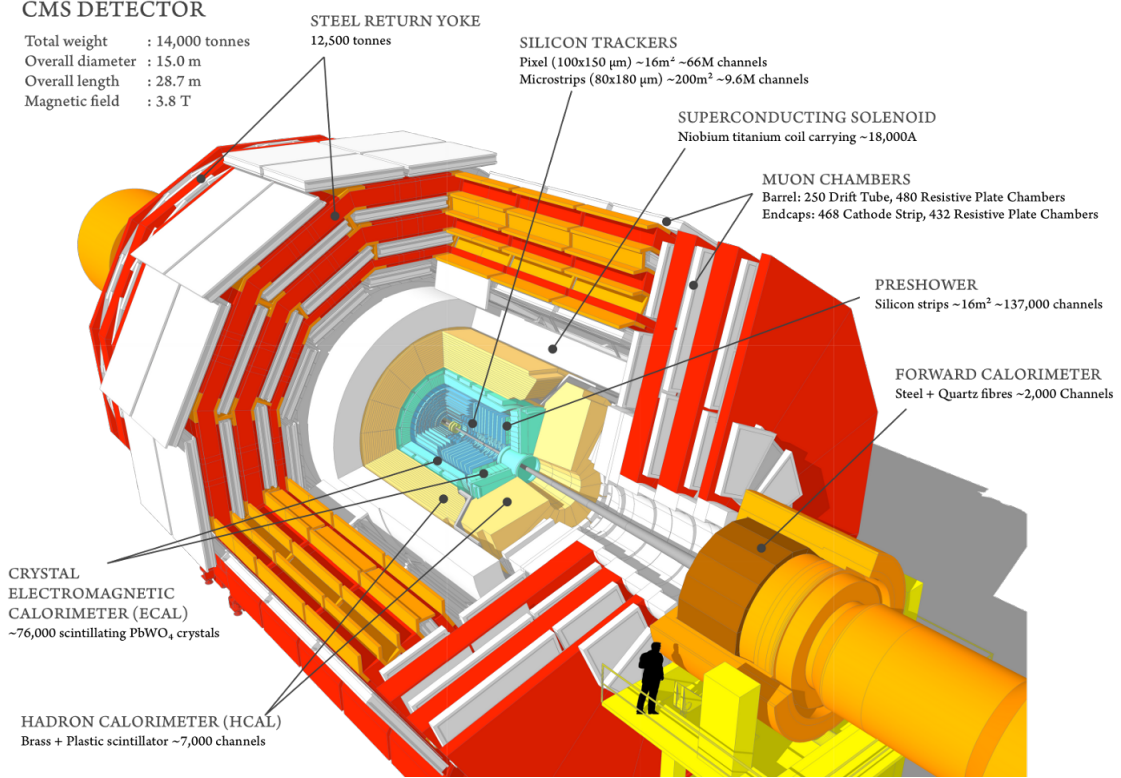


Figure 2.6: Schematic view of the CMS detector.

the two regions.

A sketch of the CMS detector is drawn in figure 2.6, highlighting the location of the various subdetectors.

2.2.2 Coordinate system

The CMS experiment uses a right-handed coordinate system, with the origin at the nominal interaction point, the x -axis pointing to the centre of the LHC, the y -axis pointing up (perpendicular to the LHC plane), and the z -axis along the anticlockwise-beam direction. One also defines r as the transverse distance: $r = \sqrt{x^2 + y^2}$. The polar angle θ is measured from the positive z -axis and the azimuthal angle ϕ is measured in the x - y plane. Instead of θ , one generally rather uses the pseudorapidity η defined as:

$$\eta \equiv -\ln \tan \frac{\theta}{2} \quad (2.4)$$

This variable is an approximation for relativistic particles (i.e. particles for which $E \gg m$) of the rapidity y :

$$y \equiv \frac{1}{2} \ln \left(\frac{E + p_z}{E - p_z} \right) \quad (2.5)$$

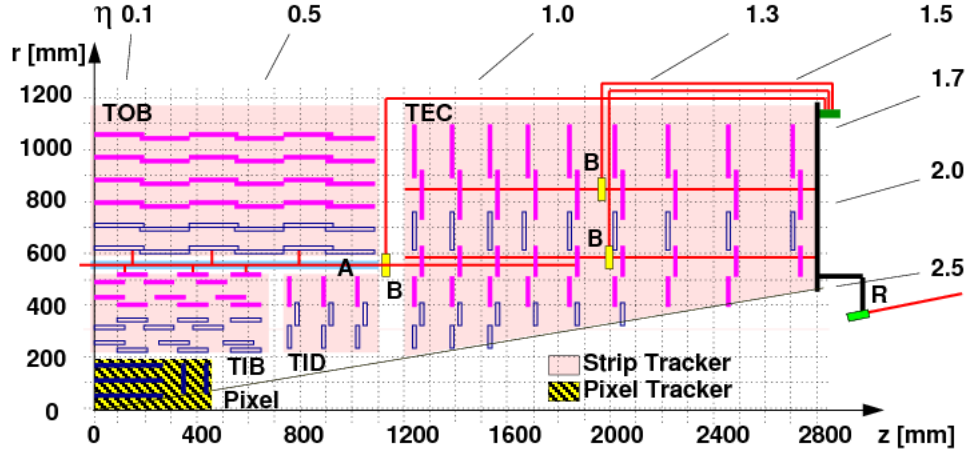


Figure 2.7: Geometry of the tracking system of CMS. Its inner part consists of a pixel detector, surrounded by a silicon strip detector divided in various regions (TIB, TOB, TEC, TID). The figure also shows the laser alignment system: the paths of the laser rays (R), the alignment tubes (A) and the beam splitters (B).

2.2.3 Magnet

The magnet used by CMS is a superconducting solenoid with an inner radius of 5.9 meters and a length of 12.9 meters. It is constituted of 2168 turns traveled by a 19.5 kA current. This choice was mainly driven by the targeted precision on the transverse momentum measurement of very high energy muons ($\approx 10\%$ for 1 TeV muons). The dimensions of the magnet allows it to embed both the tracking and calorimeter systems and to ensure a good momentum resolution in the forward region as well.

2.2.4 The inner tracking system

In order to conciliate the targeted physics performance (1.5% resolution for muons with a transverse momentum of 100 GeV) with experimental conditions (thousands of particles produced at each bunch crossing of the LHC beams), the CMS inner tracking system uses high granularity silicon detectors. As of today, it is the largest silicon tracker ever built. Its inner part is a pixel detector composed of three layers in the barrel and two in the endcaps. The pixels are surrounded by a silicon strip detector. A system of laser beams is used to monitor the position of some modules of the strip tracker. An overview of the inner tracking system is shown in figure 2.7. The tracker geometry coverage extends to $|\eta| = 2.5$.

2.2.4.1 The pixel detector

The pixel detector is made of 66 million pixels, each of them having a size of $100 \times 150 \mu\text{m}^2$ in the $r - \phi$ and z directions, allowing it to track the paths of particles emerging from the collision with extreme accuracy. The pixel cells are arranged in three cylindric layers with a radius of 4, 7 and 11 cm (pixel barrel) and in two disks in the forward regions

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at $|z| = 35$ and 47 cm (pixel endcaps). The spatial resolution is around $10\text{ }\mu\text{m}$ for the $r - \phi$ measurement and $20\text{ }\mu\text{m}$ for the z measurement. The typical charged particle flux for proton-proton collision in the first pixel layer is $10^8\text{ cm}^{-2}\text{s}^{-1}$, leading to an occupancy rate of 10^{-4} per pixel per bunch crossing.

2.2.4.2 The silicon strip tracker

The silicon strip tracker is made of 10 layers of microstrip detectors. It is divided in 4 categories:

- The Tracker Inner Barrel (TIB), consisting of the first 4 layers in the region $|z| < 65$ cm and $20 < r < 110$ cm.
- The Tracker Outer Barrel (TOB), consisting of the next 6 layers in the region $|z| < 110$ cm and $110 < r < 180$ cm.
- The Tracker Endcaps (TEC) is made of 9 disks in the forward region ($120\text{ cm} < |z| < 180\text{ cm}$ and $20 < r < 180\text{ cm}$).
- Tracker Inner Disks (TID) is composed of 3 disks that fill the gap between the TIB and the TEC.

The strips are aligned in the z direction and can therefore only measure the $r - \phi$ coordinate. However, some of them (in blue in figure 2.7) are made with double sided modules in which the second side is rotated by 100 mrad , allowing one to measure both $r - \phi$ and $r - z$ coordinates, with a resolution of $20\text{-}50\text{ }\mu\text{m}$ and $200\text{-}500\text{ }\mu\text{m}$ in the $r - \phi$ and $r - z$ direction respectively. A charged particle emitted at the interaction point and flying through the complete tracker will generate between 10 and 15 hits in the tracker, depending on the particle η .

2.2.4.3 Tracker material budget

A perfect tracking system should record the particles position without affecting their energy. However, this ideal case is far from the experimental reality as can be observed in figure 2.8, which shows the depth of the tracker, in terms of the number of radiation lengths and interaction lengths¹. There is thus a non negligible probability for the particle to interact in one of the tracker layers. This affects not only charged particles that can radiate a photon but also neutral particles: in CMS, almost half of the photons will convert in a e^+e^- pair in the tracker.

¹ The radiation length X_0 is a characteristic of a material giving the distance over which a charged particle loses in average 63% ($1 - 1/e$) of its energy through electromagnetic interaction. It is also related to the mean distance over which a photon will convert into a e^+e^- pair, Y_0 : $Y_0 = 7/9X_0$. Similarly, one also defines the interaction length λ_I as the distance over which a hadron loses in average 63% of its energy through strong interaction.

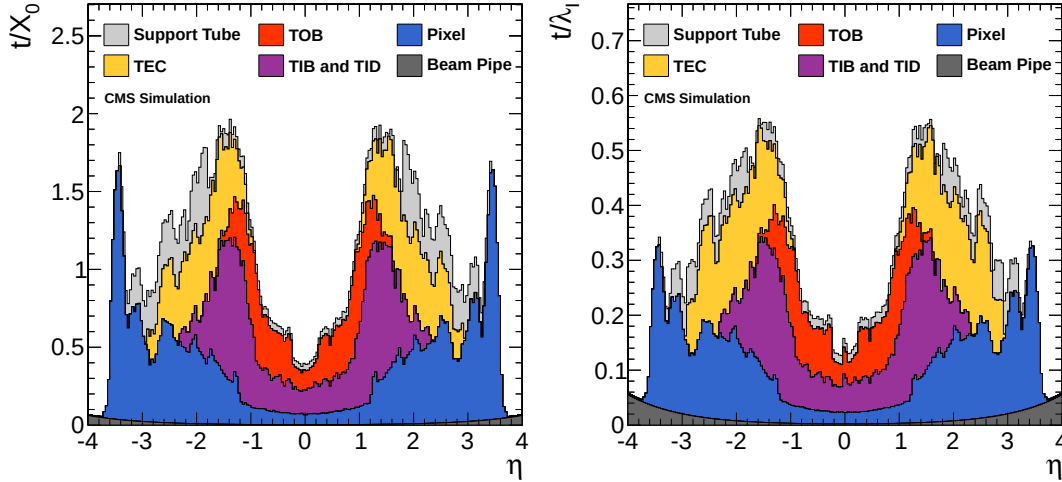


Figure 2.8: Material budget of the CMS tracker in units of radiation length X_0 (left) and interaction length λ_I (right) as a function of η for the latest version of the tracker geometry in the 2012 data taking [64].

2.2.4.4 Tracker alignment

Misalignment of the different tracker modules is one of the possible limiting cause of the physics performance of the detector. The tracker alignment was therefore performed using real data collected by CMS in 2011. Several millions of tracks of different kinds were used: muons from W and Z decays, tracks arising from diffractive inelastic interactions, and muons from cosmic rays. They were given as an input to a simultaneous fit performed on more than 200000 tracker parameters. The resulting statistical accuracy on the module alignment is lower than $10 \mu\text{m}$, below the tracker hit resolution [65].

2.2.5 Calorimeters

The first role of the calorimeters is to measure the energy of all the particles emitted at the interaction point (except muons and neutrinos), but they also exploit the different ways they interact with matter to help to distinguish them from each other. The CMS detector includes therefore two calorimeters. The closest to the interaction point is an electromagnetic calorimeter (ECAL) designed to accurately measure the energy of electrons or photons and whose depth corresponds to more than 20 interaction lengths but a small radiation length (around 1) and is surrounded by the hadronic calorimeter (HCAL) that measures the energy of hadrons.

2.2.5.1 The electromagnetic calorimeter

The CMS ECAL is a scintillating calorimeter: when a photon or an electron enters the ECAL, it starts showering and a cascade of electrons, positrons and photons are produced that excite the atoms in the material. Those atoms return to their ground state by

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emitting photons, that are collected by a photodetector. The amount of light deposited is proportional to the energy of the incident particle offering a way to measure it.

The CMS ECAL is made of 75,848 lead tungstate (PbWO_4) crystals split in a central part, the barrel ($|\eta| < 1.479$), and two endcaps ($1.479 < |\eta| < 3.0$). Each crystal weighs 1.5 kg, is 20 cm long, and has a transverse inner size of $2.2 \times 2.2 \text{ cm}^2$ in the barrel and $3 \times 3 \text{ cm}^2$ in the endcaps. The crystals were built by two producers in Russia (BTCF) and in China (SIC). The scintillation light emitted is blue-green with a maximum emission between 420 and 430 nm. Lead tungstate was chosen for several reasons: it possesses a short radiation length ($X_0 = 0.89 \text{ cm}$) and a Molière radius, defined as the radius of a cylinder containing on average 90% of the shower energy deposition, of 2.2 cm only, together with a fast response (80% of the light is emitted within 25 ns). The scintillation light is collected by avalanche photodiodes (APD) in the barrel while vacuum phototriodes (VPT) are used in the endcaps because of the higher radiation rate. The signal is then pre-amplified, shaped and amplified by three parallel amplifiers with gains of 1, 6 and 12. The maximal signal that can be recorded by a APD or a VPT corresponds to an energy deposit in the crystal of 1.7 or 3.0 TeV respectively. The signal is finally digitalized with an Analog-To-Digital Converter (ADC).

The arrangement of the ECAL crystals depends on the detector region. In the barrel, they are grouped into 16 supermodules and form rings of constant η width. The crystals are therefore more and more inclined as η increases, as shown in figure 2.9. The resulting crystal size in η and ϕ is 0.0174×0.0174 . In order to avoid that a particle escapes detection by following the path defined by the small insensitive region between two supermodules, the crystals point 3° off in η and ϕ with respect to the center of the detector. In the endcaps, however, the crystals follow a $x - y$ geometry: they form rows and lines of constant x or y .

In front of the endcaps, a preshower detector is placed. This detector is made of 2 mm wide strips, composed of two layers of lead, each one followed by a silicon sensor. This fine granularity allows it to spatially separate the two photons coming from the decay of a neutral pion, that are usually too collimated to be resolved by the ECAL endcaps. The two hits recorded in the preshower also provides a measurement of the photon direction. The preshower depth corresponds to around 3 radiation lengths.

The ECAL energy resolution depends on the energy of the incident photon or electron through the following formula:

$$\frac{\sigma_E}{E} = \frac{N}{E(\text{GeV})} \oplus \frac{S}{\sqrt{E(\text{GeV})}} \oplus C \quad (2.6)$$

where $\oplus$ denotes the quadratic sum and the three terms correspond to the noise, stochastic and constant contributions respectively. The noise term corresponds to a 40 MeV electronic noise per channel. The stochastic term is due to the variation of the number of photoelectrons and fluctuations in the gain process. The constant term finally arises from the non uniformity and stability of the responses of the various channels, the energy leakage at the back of the calorimeter, and the longitudinal non uniformity of the light collection. In the barrel, the values of N , S and C have been measured in test beams in 2004 to be 0.128, 0.028 and 0.003 [66]. Those results were obtained using the energy

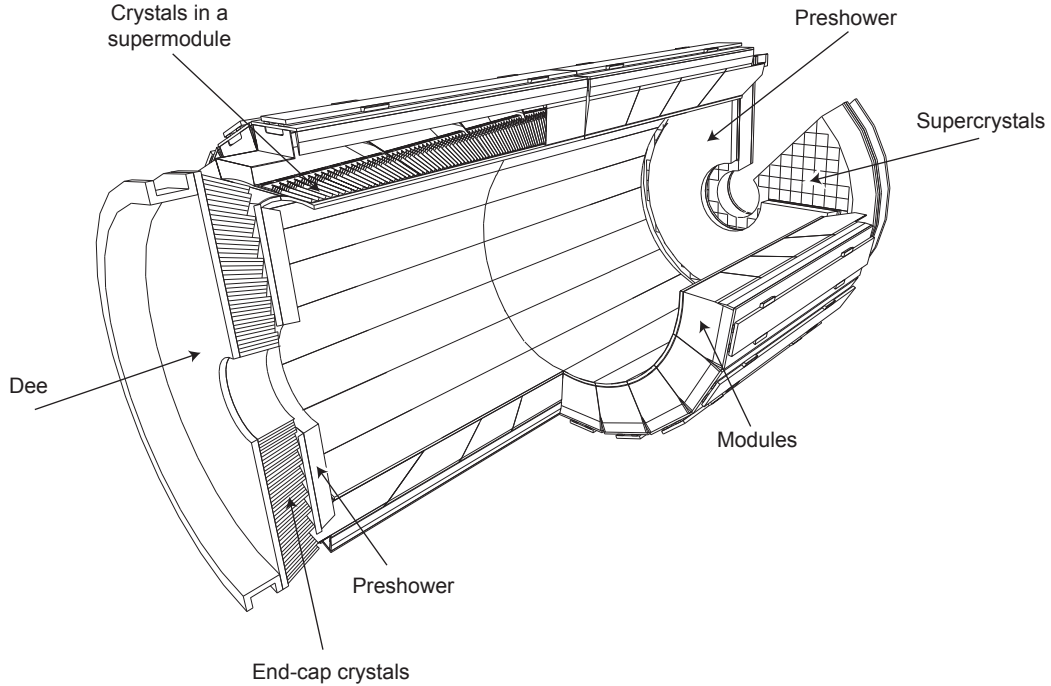


Figure 2.9: The CMS electromagnetic calorimeter and the ECAL preshower.

collected in a 3×3 matrix of crystals, for electrons with an impact point smaller than 4 mm with respect to the center of the matrix in both η and ϕ directions .

Since the ECAL plays a central role in the electron reconstruction, the following paragraphs details how it operated during the 2011 and 2012 data takings, how the energy was measured and how the detector was calibrated [67, 68].

ECAL operation

During 2011 and 2012, various operation parameters of the ECAL were monitored continuously. The crystal light yield and the APD gain show a sizeable dependency on the ECAL temperature ($-2\%/^{\circ}\text{C}$ for each), requiring a temperature stability of 0.05°C in the barrel and 0.1°C in the endcaps in order to contribute to the C term in equation 2.6 by less than 0.002. The gain variation of the APD is moreover also sensitive to the high voltage applied and its fluctuations needed to be maintained below 65 mV.

There are also some parameters that strongly varied with time. The electronic noise per crystal in the barrel was for instance measured to raise from 42 MeV at the beginning of 2011 to almost 55 MeV at the end of 2012. In the endcaps, at very high η , it evolved from 150 MeV to 600 MeV. Another time-dependent parameter is the crystal transparency, which suffers from irradiation that rises with increasing instantaneous luminosity. The crystal transparency however recovers slowly through spontaneous annealing during quiet periods. Laser light with a wavelength of 440 or 447 nm was therefore used to monitor the crystal responses every 20 to 30 minutes. The resulting relative response versus time is shown in figure 2.10. The change observed in the ECAL channels is up to 6% in the

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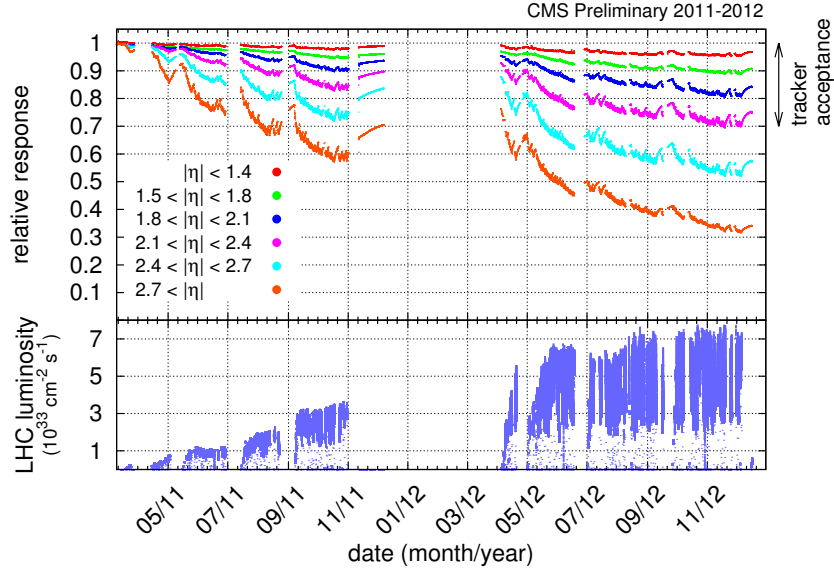


Figure 2.10: Relative response to laser light (440 nm in 2011 and 447 nm in 2012) measured by the ECAL laser monitoring system, averaged over all crystals in bins of pseudorapidity, for the 2011 and 2012 data taking periods. The bottom plot shows the instantaneous LHC luminosity delivered during this time period [68].

barrel and reaches 70% in the most forward crystals of the endcaps.

Energy measurement and calibration of photons or electrons

Since the transverse size of an ECAL crystal corresponds to about one Molière radius, an electron or a photon will leave its energy in more than one crystal. Clusters of adjacent crystals carrying a sizeable fraction of the electron or photon energy are therefore formed. In addition, the fact that electrons can emit bremsstrahlung photons and that photons can convert in the tracker implies that the energy deposit corresponding to the primary particle can actually be split into various clusters due to the electron bending in the magnetic field, and supercluster algorithms have been developed to group all the individual clusters. They are detailed in section 4.1.1.

The photon or electron energy, $E_{e,\gamma}$, is estimated starting from the signal amplitude in ADC counts, A_i , in each crystal belonging to the supercluster. It is given by:

$$E_{e,\gamma} = F_{e,\gamma} \cdot \left(G \cdot \sum_i S_i(t) \cdot C_i \cdot A_i + E_{ES} \right) \quad (2.7)$$

where G is a global calibration constant (different in the barrel and in the endcaps) converting the ADC counts in GeV and only measured in a limited region of the ECAL, C_i are intercalibration coefficients, $S_i(t)$ are correction factors to account for the time dependency crystal transparency and E_{ES} is the preshower energy for superclusters reconstructed in the endcaps. Finally, $F_{e,\gamma}$ represents a correction to the supercluster energy that takes into account η and ϕ dependent effects such as those due to the material budget

in front of the supercluster. Since electron and photon showers are not exactly identical, this factor is estimated separately for the two cases.

The $S_i(t)$ factor is estimated using the ECAL light monitoring (LM) system: the crystal response to laser light, $R_i(t)$, is first measured. Since the spectral composition of this light differs from the scintillation light of the crystal, $S_i(t)$ and $R_i(t)$ are not equal and the following relation is assumed:

$$\frac{S_i(t)}{S_{i0}} = \left(\frac{R_i(t)}{R_{i0}} \right)^\alpha \quad (2.8)$$

where S_{i0} and R_{i0} are the initial response to scintillation and laser light respectively. The values of α were measured in test beam for crystals from the two producers and found to be 1.52 for BTCP and 1.00 for SIC. Three methods were used to tune and validate the response corrections:

- A fit to the diphoton mass distribution was performed in the region of the η meson mass and the stability of the η mass measurement was studied as a function of time. The large event statistics allows one to quickly validate the corrections on short periods of time (less than 1h).
- The time evolution of the ratio of the energy over momentum (E/p) measured on isolated electrons coming from W decay was also studied. Figure 2.11 shows the E/p scale factor, measured in 2012 before and after the response corrections, that needs to be applied to match the E/p distribution measured in the full 2012 dataset with response corrections applied. The fact that the scale factor is 1 when corrections are applied indicates that the correction factor measured from laser light is properly estimated.
- Events where a Z boson decays into a dielectron pair were used to tune the α parameter for BCTP crystals: the value of α minimizing the Z mass resolution while maintaining a stable E/p scale factor was taken.

The intercalibration constants C_i relate the global calibration G to its exact value in each crystal. They are estimated by combining three measurements:

- The ϕ -symmetry method exploits the fact that the average energy deposit in each crystal should be the same for a given η . The energy deposits measured in generic inelastic proton-proton scattering events are therefore used to derive intercalibration constants for crystal belonging to the same ϕ ring at constant η . The precision of the measurement is limited by the uncertainty on the azimuthal inhomogeneities in the detector.
- The decays of π_0 and η mesons into two photons offers another way to intercalibrate the crystals using the $\gamma\gamma$ invariant mass. For crystal i , the invariant mass distribution of photon pairs where one of the photons is reconstructed in the center of a this crystal, is fitted to the sum of a Gaussian distribution, modeling the signal, and a fourth order polynomial modeling the background, allowing one to extract a measurement m_i of the π_0 or η mass. The intercalibrations are fixed through an

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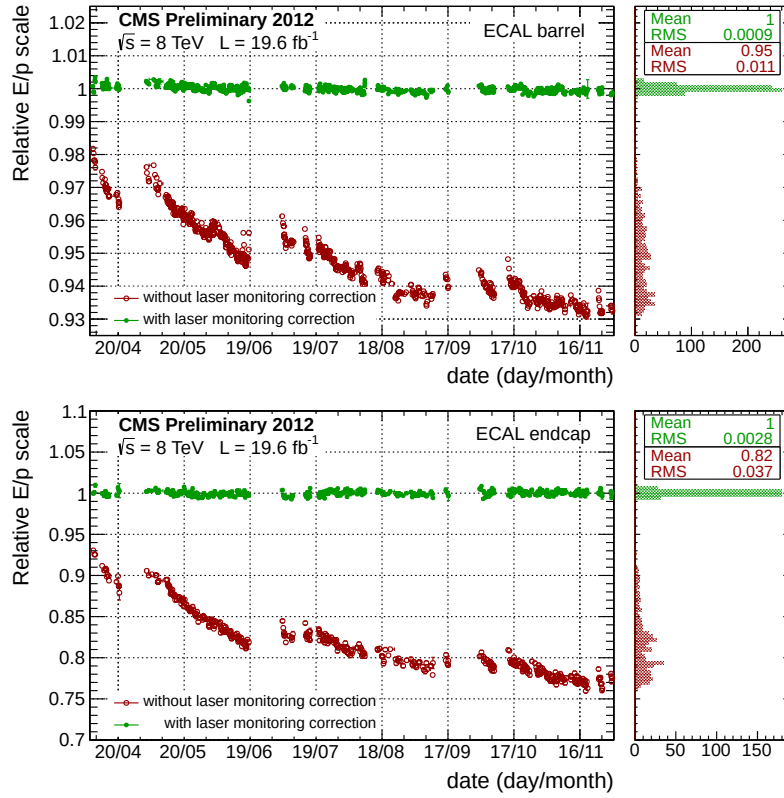


Figure 2.11: Relative scale on the E/p ratio that needs to be applied to match the corrected E/p distribution measured in the full 2012 dataset as a function of time before (red) and after (green) response corrections for electrons reconstructed in the ECAL barrel (top) and endcaps (bottom). The right part of the plots presents a histogram of the ratios obtained and can be used to estimate the uncertainty on the corrections [68].

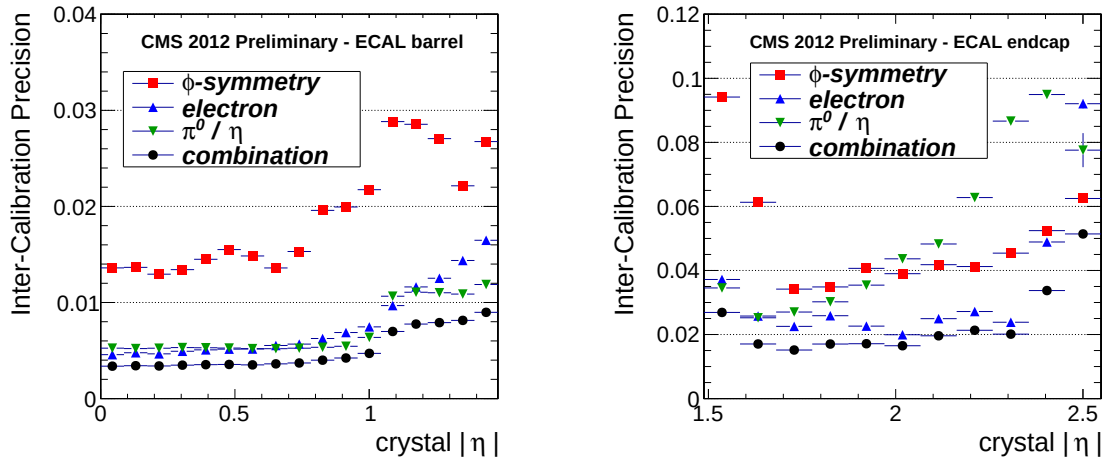


Figure 2.12: Precision on the single crystal calibration as a function of the crystal η using the ϕ -symmetry method (red), π_0 and η decays into two photons (green) and electrons from W or Z decays (blue) obtained using data recorded in 2012. The precision on the combined measurement is also presented (black). Results are shown separately for ECAL barrel (left) and endcaps (right) crystals [68].

iterative procedure until all the m_i measurements converge. The uncertainty on this method arises from the precision of the fit.

- Similarly to the time dependent response corrections, the E/p ratio measured on electrons from W and Z boson decays offers another measurement of the intercalibration constants. An iterative calibration procedure was used in each ϕ ring at constant η : the E/p distribution for electrons in each crystal were scaled until they match the global E/p distribution for all the crystals in the ring. A η calibration was then performed. Since the E/p distribution depends on η due to the change of the tracker material budget that impacts the energy and momentum resolution, reference distributions were generated from simulation in various η regions and scaled to the data. In this case, the number of electrons per crystal is the limiting factor to the precision that can be reached.

The intercalibration precision for each measurement and for their combination are displayed in figure 2.12.

The energy deposited in the preshower, E_{ES} , represents between 6 and 8% of the total energy of an electron coming from the decay of a Z boson. This fraction decreases with the electron or photon energy. It is given by the following formula:

$$E_{ES} = G_{ES} (E_{ES}^{clus1} + \alpha_{ES} \cdot E_{ES}^{clus2}) \quad (2.9)$$

where E_{ES}^{clus1} and E_{ES}^{clus2} are the energy deposits expressed in number of minimum ionizing particles (MIPs) in each of the preshower layers. Those factors were calibrated using charged pions and muons with momentum greater than 1 GeV recorded in the data. The coefficient α_{ES} represents the weight of the second layer with respect to the first layer.

2. THE EXPERIMENTAL SETUP

It was measured in test beams to be 0.7. Finally, the G_{ES} factor converts the number of MIPs in an energy measurement in GeV. It is estimated from a linear fit of the ECAL cluster energy as a function of $(E_{ES}^{clus1} + \alpha_{ES} \cdot E_{ES}^{clus2})$ using electrons from W decays.

The energy corrections $F_{e,\gamma}$ aim to provide a customized correction to the supercluster energy in order to match as closely as possible the initial electron or photon energy. A multivariate regression technique trained on simulated samples of photon + jets or Z and W boson decays events is therefore used to estimate the ratio of the electron or photon energy at the generator level over the supercluster energy. It takes as inputs the supercluster η and ϕ and several other variables that provide local information concerning the various clusters inside the supercluster, the shower shape and the likelihood of photon conversion,... The number of extra vertices in the event is also included. This approach allows one to accurately estimate the amount of energy lost in insensitive regions of the detector or recorded outside the supercluster and to subtract energy deposits from pile-up events.

The resulting average supercluster energy corrections for electrons from W boson decays recorded during the 2011 data taking are presented in figure 2.13 (left). Electrons are split into two categories depending whether the ratio of the energy recorded in the 3×3 crystal matrix centered on the highest energy crystal, R9, is greater than 94% of the total supercluster energy or not. In the first case, electrons have usually radiated only a small fraction of their energy through interaction with the material in front of the ECAL and their correction factor F_e provides a good description of the corresponding factor for unconverted photons F_γ . Figure 2.13 (right) illustrates the fact that the η dependency of the $F_{e,\gamma}$ correction is related to the material budget in front of the ECAL.

The absolute calibration of the ECAL, G , is computed in a region of the detector where the energy corrections and uncertainties are minimal. In the barrel, it consists of the 150 crystals with $|\eta| < 0.35$ and distant by at least 5 crystals from the border of a module. In the endcaps, it is defined as all the crystals with $1.7 < |\eta| < 2.1$. G is first estimated from unconverted photons in a simulated sample such that the photon energy is equal to the energy measured in a 5×5 crystal matrix. The value obtained from simulation is then corrected using electrons pairs from Z decays: the Z mass is measured from a fit of the dielectron invariant mass distribution in the simulation and in the data. The scale factor is then defined as $(1 + \Delta m)/M_Z$ where Δm is the mass difference between the two measurements and $M_Z = 91.188$ GeV. In order to cross check this result with photons rather than electrons in the data, Z boson decays into two muons where one of them radiates a photon is also used. The procedure is similar: the invariant mass distribution of the 2 muons and the photon is fitted and an energy scale is extracted. Compatible corrections are found with the two methods. The uncertainty of the absolute energy calibration is mainly due to the uncertainty on $F_{e,\gamma}$ and was determined to be around 0.4% in the barrel and 0.8% in the endcaps.

The effect of the various corrections discussed above on the dielectron invariant mass distribution for Z decays events is illustrated in figure 2.14.

ECAL resolution

The energy resolution is primarily estimated from simulation. It is corrected by comparing the experimental width of the Z peak in the data and in simulation. For electrons

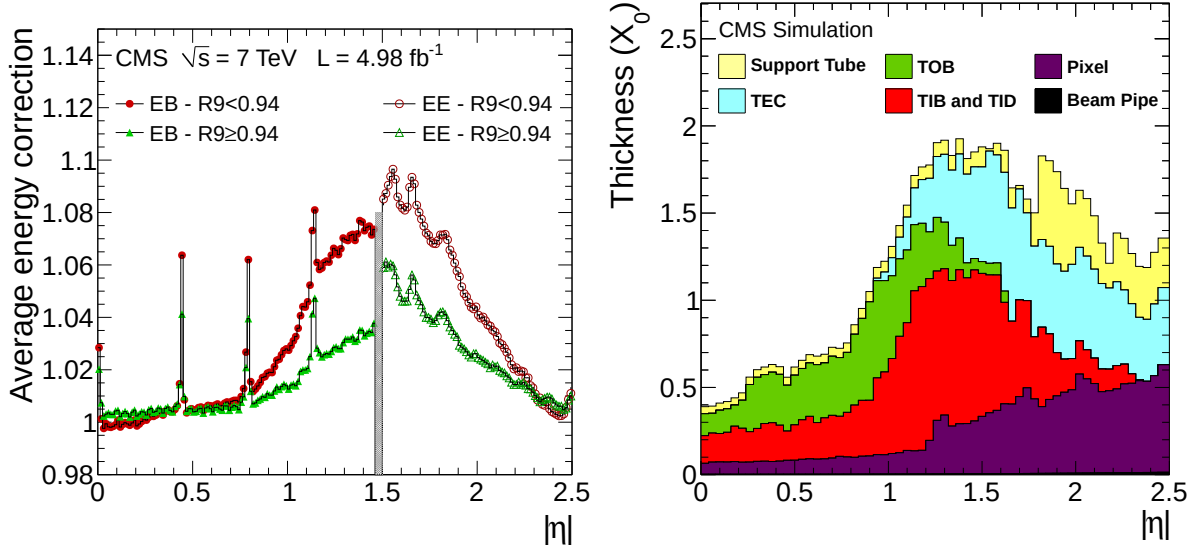


Figure 2.13: Left: average value of the F_e correction as a function of the electron η measured on electrons from W decays recorded in 2011 data for two R9 regions. The spikes observed in the ECAL barrel are due to the energy leakage between the supermodules. Right: material budget in front of the ECAL as a function of η [67].

produced from Z decays and photons from a 125 GeV scalar boson, the resolution is comparable and ranges from 1 to 5% depending on the η region considered [67]. Those values are larger than the ones obtained from equation 2.6 because the possible interaction of electrons and photons within the tracker enlarges the shower spread in the ϕ direction in the presence of a magnetic field. The evolution of the resolution with energy is shown in figure 2.15 where the energy resolution of the ECAL for a prompt and isolated electron recorded in the barrel is compared to the one obtained using the track information from the tracker. The ECAL resolution at higher energy (up to a few TeV) is presented in section 6.2. For large energy values (> 1.7 TeV in the barrel and > 3.0 TeV in the endcaps), the ECAL crystal saturate. A method to estimate the energy in such crystals using the energy deposits in the surrounding crystals is described in section 6.2.1.

2.2.5.2 The hadronic calorimeter

The hadronic calorimeter (HCAL) of CMS is designed to measure the energy of the hadrons jets produced in the proton-proton interaction. It also plays a central role in the identification of neutrinos and exotic particles that do not interact with the detector and whose signature is a non zero transverse energy balance. Since the HCAL is placed inside the magnet, materials with short interaction lengths need to be used. A good hermiticity is also requested in order to accurately measure the transverse missing energy.

The CMS HCAL is a sampling calorimeter consisting of a succession of layers of absorbers (mainly brass) and scintillator material.

The central part (HB) which extends from $|\eta| = 0$ to $|\eta| = 1.3$ is made of 16 absorber

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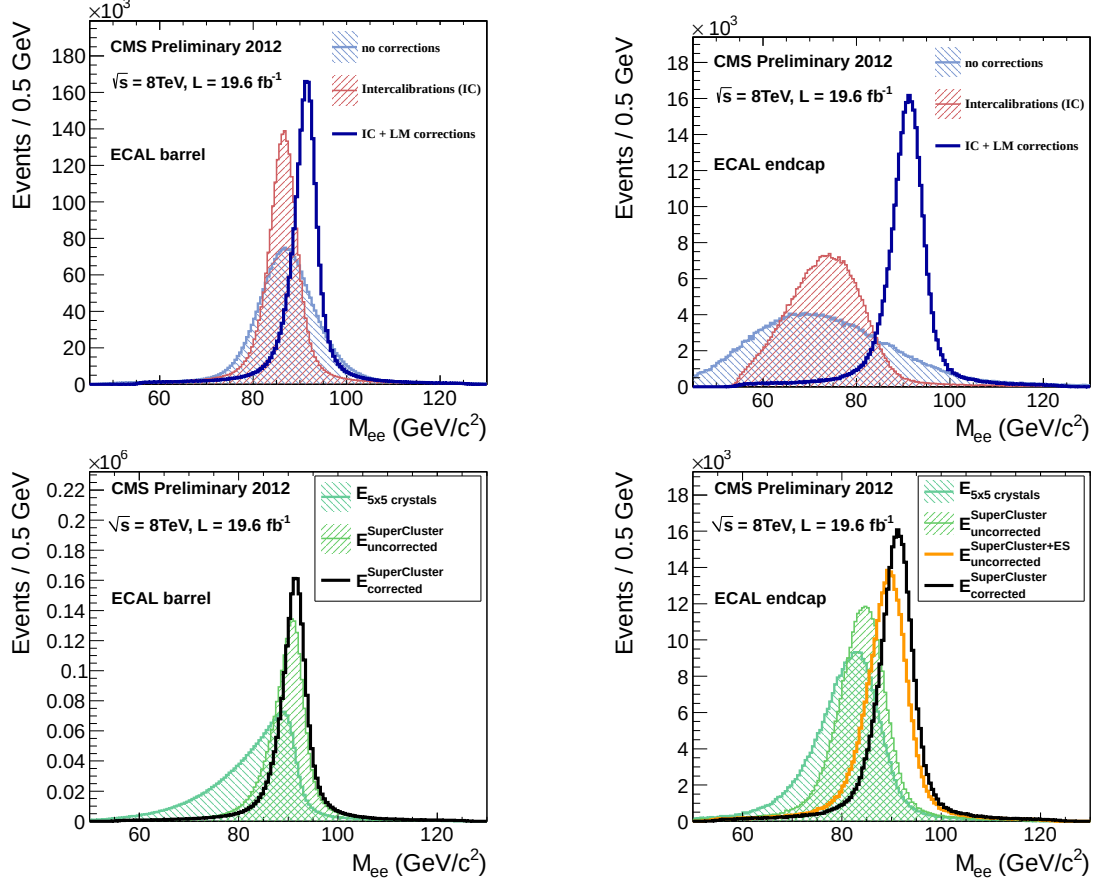


Figure 2.14: Top: reconstructed dielectron invariant mass for electrons from Z decays events computed before the $S_i(t)$ and C_i corrections are applied (blue), when only the C_i correction is applied (ref) and when both corrections are applied (black). In all the cases, the supercluster energy F_e correction was applied and the preshower energy is included in the endcaps. Bottom: reconstructed dielectron invariant mass for electrons from Z decays events computed using the energy measured in a 5×5 matrix centered on the highest energy crystal (dark green), the supercluster energy before the F_e correction is applied (light green) or after correction and including the preshower energy in the endcaps (black). The orange curve in the endcaps corresponds to an electron energy estimated as the sum of the uncorrected supercluster energy and the preshower energy. Results are shown separately for events where both electrons are reconstructed in the barrel (left) or in the endcaps (right) [68].

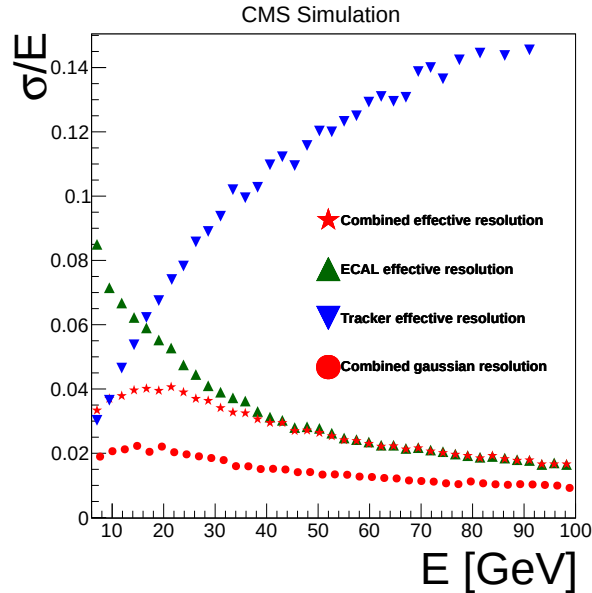


Figure 2.15: Expected energy resolution for prompt and isolated electrons emitted in the barrel as a function of the electron energy, obtained using the ECAL (green) or tracker (blue) information, or a combination of both (red). The resolution is evaluated as half the minimal width that contains 68.3% of the reconstructed energy or momentum distribution (effective resolution) and using a Gaussian fit of the core of the momentum distribution [8].

plates. The innermost and outermost ones are made of stainless steel whereas the others consist of brass. The interleaved plastic scintillators are segmented into 32 regions in η times 36 regions in ϕ , leading to 1152 towers 0.087×0.087 wide in $(\Delta\eta, \Delta\phi)$. On each side of the barrel, the two towers closest to the endcap transition region are also longitudinally segmented in two depths. Because of the barrel geometry, the number of interaction lengths of the HCAL varies between 5 ($|\eta| = 0$) and 10 ($|\eta| = 1.3$). This is sometimes too small to contain the full hadronic shower. In order to recover the end of the shower, the HCAL barrel is completed by an outer calorimeter (HO) surrounding the magnet, raising the total depth of calorimeter system to at least 12 interaction lengths.

The HCAL endcaps (HE) extend from $|\eta| = 1.3$ to $|\eta| = 3.0$. Its $\Delta\eta \times \Delta\phi$ granularity is 0.087×0.087 for $|\eta| < 1.6$ and around 0.17×0.17 for $|\eta| > 1.6$. As in the barrel the absorber is made of brass plates. The total depth of the endcap calorimetry is about 10 interaction lengths. The HCAL endcaps are longitudinally segmented into two depths for $|\eta| < 2.5$ and into three depths at higher $|\eta|$.

A forward calorimeter (HF) is also installed at 11.2 m from the interaction point that covers the $|\eta|$ range from 3.0 to 5.2. Due to the extreme particle fluxes in that region, a radiation hard technology had to be used. The choice fell on Cherenkov light detector made of quartz fibres. In addition to ensuring a good hermiticity of the HCAL, this detector is used for the measurement of the instantaneous luminosity.

The position of the various HCAL components can be seen in figure 2.16.

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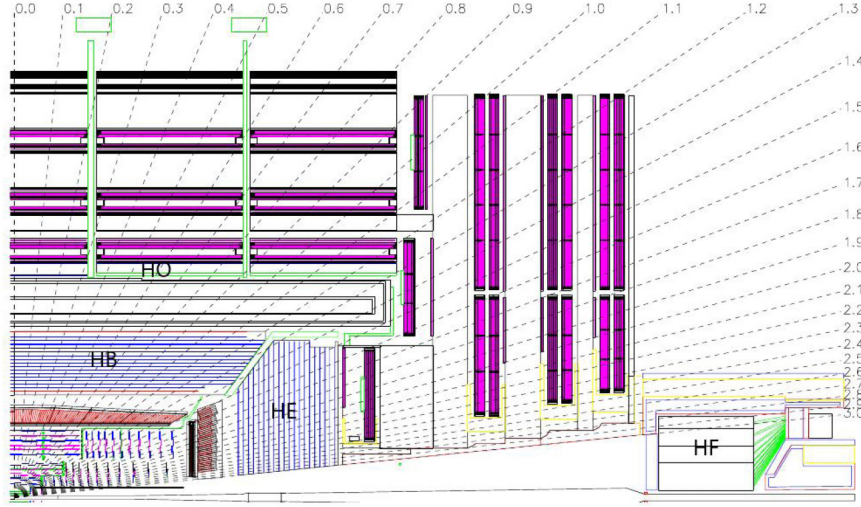


Figure 2.16: Schematic view of a slice of a quarter of the complete CMS detector indicating the position of the various components of the hadronic calorimeter: the HCAL barrel (HB) and endcaps (HE), the outer HCAL (HO) and the forward HCAL (HF).

2.2.6 Muon chambers

One of the key components of the CMS detector is its muon chambers. The muon is the only known charged particle that can pass through the whole detector almost without interacting. The muons chambers are designed to track those particles and to measure their momentum. They are placed outside the magnet and consist of three different kinds of gaseous detectors in CMS.

In the barrel ($|\eta| < 1.2$), 4 stations of drift tubes (DT) are placed between the layers of the magnet return yoke. The three stations closest to the interaction point are divided into 8 chambers measuring the $r - \phi$ coordinate and 4 others measuring the z coordinate. The last station only measures the $r - \phi$ coordinate. The resolution provided by each station is around $100 \mu\text{m}$ in position and 1 mrad in direction.

In the forward regions ($0.9 < |\eta| < 2.4$), the magnet field is too intense and the neutron-induced background becomes too high to be sustainable for drift tubes. Cathode strip chambers (CSC), consisting of multiwire proportional chambers, are therefore used. The CSC resolution is however not as good as the DT: around $200 \mu\text{m}$ in position and 10 mrad in direction.

Both the DT and CSC can be used to trigger an event containing a muon with high transverse momentum. Their time response however is comparable to the time between two consecutive bunch crossing. The muon system is therefore completed by three layers of resistive plate chambers (RPC) covering the $|\eta|$ range from 0 to 1.6. The RPC offer a much coarser spatial resolution than the other technologies but provide a time resolution of the order of 1 ns . The position of the various muon detectors are shown in figure 2.17. The muon momentum is measured by combining information from the inner tracker and muon chambers. At low transverse momentum (below 200 GeV), the resolution is completely driven by the tracker and is around 1% (resp. 2%) in the barrel (resp. in the endcaps)

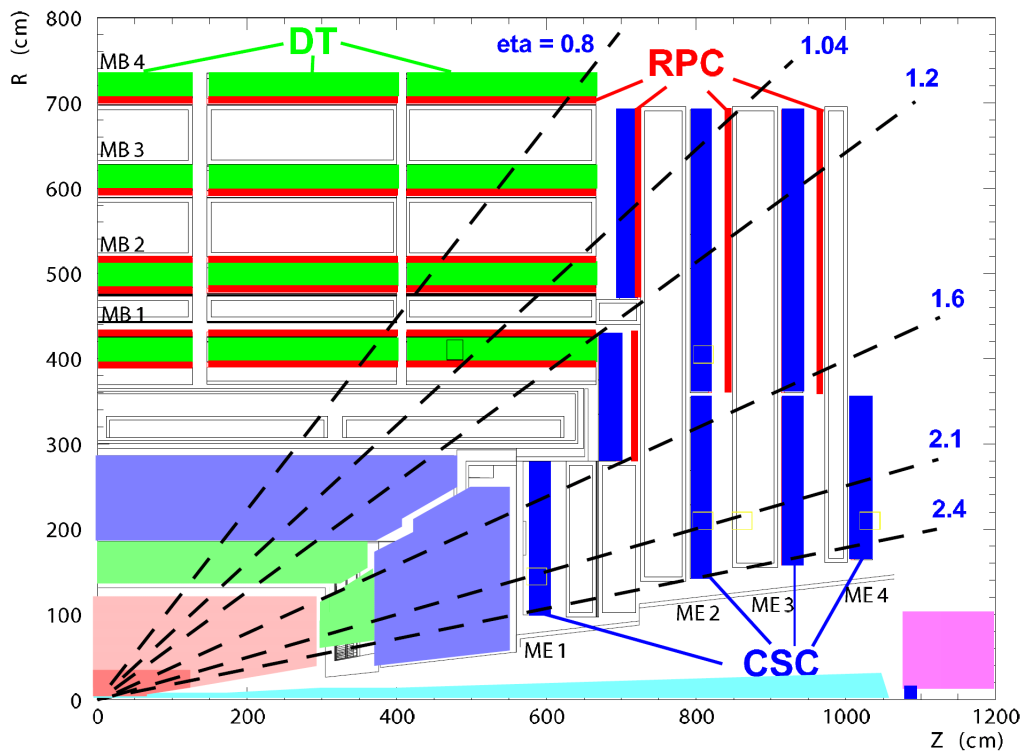


Figure 2.17: Sketch of the CMS muon system. Drift tubes (DT) and cathode strip chambers (CSC) provide an accurate position measurement. The resistive plate chambers (RPC) ensure a time resolution well below the time separating two bunch crossings.

2. THE EXPERIMENTAL SETUP

for a 50 GeV muon. At higher values, both the tracker and muon chamber contribute, leading to a resolution of 5% (resp. 10%) in the barrel (resp. in the endcaps) for a 1 TeV muon.

2.2.7 The trigger system

The LHC is designed to collide bunches of protons every 25 ns, with an average number of 20 interactions per bunch crossing. This represents an amount of data much too high to be recorded. A system of trigger is therefore used to save only the most interesting collisions, characterized by the presence of particles with high transverse momentum. The reduction of the rate is achieved in two steps called the level 1 trigger (L1) and high level trigger (HLT).

Each of them builds a set of trigger objects and the decision to keep or discard an event is based on various criteria, such as the type of those trigger objects (electrons, photons, muons...), their number, the quality of the reconstruction or their transverse energy or momentum. In order to still keep a fraction of the most frequent collisions (which is in particular needed to study pile-up), prescaled triggers are also used. The final trigger selection is a programmable menu that is made of up to 128 alternative individual conditions.

- The L1 trigger consists of programmable electronics relying on the field-programmable gate array (FPGA) and Application-specific integrated circuit (ASIC) technologies. It aims to reduce the rate from 40 MHz to 100 kHz. The L1 trigger uses only the information coming from the calorimeters and the muon chambers as its requested latency of 3.2 μ s forbids the use of the tracks whose reconstruction is too slow. The calorimeters are coarsely segmented into trigger towers that are analyzed by a Regional Calorimeter Trigger (RCT). The RCT measures the transverse energy sums, identifies regional photon or electron candidates (up to $|\eta| = 2.5$) and tau candidate (up to $|\eta| = 3.0$). Information from the different RCT are then gathered by the Global Calorimeter Trigger that establishes a set of trigger objects accross the entire detector. A same chain holds for the muon trigger: the information from the DT, the CSC and the RPC are first treated separately in a Regional Muon Trigger and then put together by the Global Muon Trigger. The rapidity coverage of the muon triggers is $|\eta| < 2.1$. Events passing the L1 trigger are given as input to the HLT trigger.
- The HLT trigger is in charge of further reducing the rate down to 100 Hz. It consists of a filter farm of around 1000 processors that can access the full read out information. The typical event size at the HLT is 1.5 MB. It is at the HLT level that the distinction between photons and electrons, who have identical signatures in the calorimeters, is performed.

The electron triggers used to collect the data analyzed in the current thesis are described in details in section 4.2.

2.3 Object reconstruction with the CMS detector

The inelastic collision of two high energy protons at the center of the CMS detector produces a large number of particles in the final state which must be properly reconstructed and identified.

They can be classified in objects with different signatures in the detector: electrons and photons (section 2.3.2), jets of hadrons coming from the hadronization of a quark or a gluon (section 2.3.3), muons (section 2.3.4) and tau leptons which are identified through their decay products (section 2.3.5). The last type of particle present in the final state is the neutrino, whose presence can be inferred from an imbalance of the transverse energy (section 2.3.6). The measurement of a track in the tracker detector for charged particles and the reconstruction of the primary vertices represent key aspects of the reconstruction of the various objects and are detailed in section 2.3.1.

2.3.1 Track and vertex reconstruction

A particle passing through a tracker layer can in general induce a signal in more than one pixel or more than one strip. The first step of the tracking procedure is the clusterization which assembles nearby tracker channels into one hit cluster. The particle position and its uncertainty is then inferred from the relative signal amplitudes in each channel.

Since the magnetic field is almost constant within the tracker, the trajectory of charged particles can be modelled as a helix parametrized by 5 parameters. The track reconstruction starts therefore from a seed consisting of three hits in the pixel detector or two hits and a constraint coming from the beam position, which provides a first estimate of the track parameters and of its covariance matrix.

New hits in the next tracker layers are then added using a Kalman filter algorithm [57, 69] which takes into account the fact that, due to multiple scattering of the particle in the tracker layers, the track parameters differ at each layer. Its principle is to iteratively propagate the current track trajectory to the next tracker layer and to look for a compatible hit in a window corresponding to the trajectory uncertainty. The track parameters and covariance matrix at the new layer are then updated by performing a least squares minimization which includes the information from the new hit. Gaussian errors on the measurement and on the noise propagation are assumed.

In general, more than one compatible hit are found at each layer and many tracks are reconstructed from a same seed. To limit their number, various criteria are required such as on the normalized χ^2 of the track trajectory, and on the number of valid and invalid hits. It can also happen that hits are used in several tracks from different seedings. A cleaning is therefore applied for tracks sharing a lot of common hits: only the best reconstructed one is kept while the others are discarded.

Finally, once all the hits of a track have been found, its parameters are refitted. The usual fit starting from the interaction point to the end of the tracker is complemented with a second fit run backward from the outermost tracker layer to the interaction point.

The vertex reconstruction [70] uses pixel hit triplets or fully reconstructed tracks, the advantage of the former being its fast computing time, allowing it to be used in the HLT trigger. Tracks with a small transverse impact parameter with respect to the beam spot

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are used to form various vertices spread along the z -axis. A fit is therefore performed to estimate the exact vertex position. The vertices are then sorted by decreasing values of the sum of the squared tracks transverse momenta. The first vertex is usually the one corresponding to the collision of interest (where a hard process, *cf.* section 3.1.2, takes place).

2.3.2 Electrons and photons

The reconstruction algorithms for electrons and photons are similar due to the similarities in the objects' signatures in the calorimeter. Their entire energy is measured in the ECAL (almost no deposit in the HCAL is expected) and electrons or photons entering the ECAL lead to a shower with a small lateral extension. The presence of a track matching the energy deposit allows one to usually disentangle them.

The fact that a photon can convert into a dielectron pair in the tracker and that electrons can radiate bremsstrahlung photons which can in turn convert implies however that in both cases the energy of the incident particle can be distributed between several energy deposits largely spread in the ϕ direction, due to the bending of the electrons in the magnetic field.

In some cases this feature also makes the electron/photon disambiguation difficult and various algorithms have been developed to identify the primary particle.

The final selection of electrons or photons includes various conditions (on the shower shape, on their isolation in the detector,...) in order to efficiently distinguish them from jets. The selection can take the form of a serie of cuts on several variables or of a cut on the output of a multivariate analysis. Various working points, depending on the energy range considered and on the desired efficiency and purity were developed. A detailed presentation of the electron reconstruction and identification, focused on the high energy case, is presented in chapter 4.

2.3.3 Jets

Jets arise from the hadronization of quarks and gluons which will be discussed in section 3.1.5. They are made of a large number of neutral and charged hadrons. In order to reconstruct them, Particle Flow objects are used. The Particle Flow algorithm [71, 72] aims to identify and reconstruct all the particles from the collision by combining optimally the information of the different subdetectors. Various objects are built: electrons, muons, photons, charged and neutral hadrons. Charged hadrons are reconstructed from the association of tracks with energy clusters whereas the clusters not associated to any track are attributed to neutral hadrons. Tracks and clusters already associated to an electron, a muon or a photon are not taken into account.

To group the various Particle Flow objects into a jet, several algorithms can be used. The main algorithm used in CMS is the so-called anti- k_t algorithm [73]. One first lists all the “distances” between any object i and any object j , d_{ij} , and the distances between

any object i and the beam, d_{iB} , defined as:

$$d_{ij} = \min\left(\frac{1}{k_{t,i}}, \frac{1}{k_{t,j}}\right) \cdot \frac{\sqrt{(y_i - y_j)^2 + (\phi_i - \phi_j)^2}}{R} \quad (2.10)$$

$$d_{iB} = \frac{1}{k_{t,i}} \quad (2.11)$$

where $k_{t,i}$ and $k_{t,j}$ are the object transverse momentum, y_i and y_j their rapidity and ϕ_i , ϕ_j their azimuthal angle. R is a threshold value taken to be 0.5. If the smallest value of the distances is one of the d_{iB} , the corresponding object forms a jet and is removed from the list. If it is a d_{ij} , objects i and j are grouped to form a new object with a transverse momentum equal to $k_{t,i} + k_{t,j}$ and j . The process is repeated until all the objects have been assigned to a jet.

This algorithm offers the advantage that the shape of the jet is usually close to a cone with an area of πR^2 and does not depend on the low p_t objects while still including them, as opposed to other algorithms.

The raw jet energy is subject to several corrections. First, the energy coming from pile-up events must be subtracted. It is estimated from the contribution of charged hadrons associated to vertices other than the first vertex, multiplied by a factor estimated from simulations which adds the contribution from neutral hadrons. The calorimeters response is in addition dependent on the energy and the pseudorapidity, and energy corrections, calculated from simulations, are therefore applied. Finally an extra correction is applied to take into account the different energy scale and resolution measured in data and in the simulation.

The final energy resolution for a jet of 100 GeV of p_t is around 10%.

Except the top quark, which decays before hadronizing, all the quarks will generate a jet. Techniques were however developed to distinguish b quarks from lighter ones. They are based on the fact that the hadrons containing a b quark are not stable but have a sizeable lifetime and that their decay sometimes produces leptons. The presence of a displaced vertex, of tracks with a large impact parameter with respect to the beam position and of leptons inside the jet can therefore be used to separate b-jets from lighter flavour jets.

2.3.4 Muons

Because of their mass and the fact that they are insensitive to the strong interaction, the muon energy loss in matter is extremely reduced and a muon emitted in the center of the detector leaves only a small fraction of its energy in the calorimeters. Its long lifetime (2.2 μ s) in addition ensures that the probability for a muon with an energy in the GeV range or more to decay in the detector is negligible.

The muon reconstruction is therefore only based on the tracker and muon chambers information. Tracks are first reconstructed separately in each subdetector. Two reconstructions are then used. The global muon reconstruction consists in matching a track in the tracker and a track in the muon chambers by comparing their parameters propagated to a common surface using a Kalman filter. A fit on both tracks is then performed to

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Variable	Condition	Comment
Tracker & Global muon	true	The muon must pass both reconstructions.
p_t	$> 45 \text{ GeV}$	Needed to reach the plateau of the trigger.
$ \eta $	< 2.4	Geometrical limit of the muon system.
$\delta p_t/p_t$	< 0.3	Relative uncertainty on the p_t .
# hits in tracker layers	≥ 6	
# hits in pixel detector	≥ 1	
# hits in muon system	≥ 1	
d_{xy}	$\leq 0.2 \text{ cm}$	Transverse impact parameter with respect to the vertex with which the muon candidate is associated.
# of muon segments	≥ 2	The tracker muon must match segments in at least two muon stations.
Tracker Isolation	$< 0.1 \times p_t$	The sum of the p_t of all the tracks in a cone around the muon should be less than 10% of the muon p_t .

Table 2.2: The various conditions forming the high energy muon selection.

extract new track parameters. Another reconstruction only requires a track in the tracker and looks for a track segment in the muon system. The two objects are then grouped if some matching conditions are satisfied. This second method allows one to recover some efficiency for muon with momenta of 5 GeV or less.

On top of the muon reconstruction, various muon selections are used in CMS. As an example, the high energy muon selection used in the search for heavy resonance decaying into a dimuon pair is described in table 2.2. The probability for a dimuon pair to be reconstructed and that both muons pass the high energy muon selection (“acceptance times efficiency”) is shown in figure 2.18 as a function of the dimuon invariant mass. Since the muon momentum only relies on the track information, its resolution worsens with increasing momentum: the typical mass resolution for 1 TeV dimuon pair is 4% and rises to 10% for a mass of 3 TeV.

2.3.5 Taus

Tau leptons are characterized by a short lifetime ($2.9 \cdot 10^{-13} \text{ s}$) and their decay, which produces undetectable neutrinos, occurs inside the detector. They are therefore reconstructed from their visible decay products. In the case of a leptonic decay (about one third of the cases), the electron or the muon is the only reconstructed particle and it is reconstructed as explained in the previous sections.

When the tau decays hadronically (two thirds of the cases), a jet is produced. This jet shows however specific features. It is characterized by a small track multiplicity: the hadronic decay products of a tau contain only one or three charged particles, mostly pions, which is much less than the usual track multiplicity for gluon or quark jets. The neutral part consists mostly of neutral pions which decay into two photons.

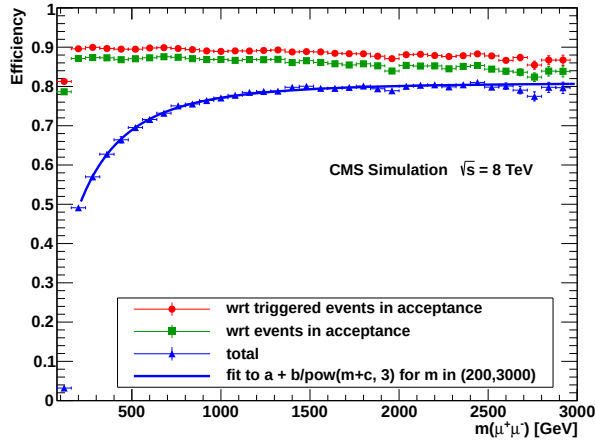


Figure 2.18: The probability for a dimuon pair to be reconstructed and that both muons pass the high energy muon selection (“acceptance times efficiency”) as function of the invariant mass of the pair. The two muons are also required to carry an opposite charge. The results are obtained from Drell-Yan simulated events at $\sqrt{s} = 8$ TeV.

In order to tag a jet as coming from a tau decay, algorithms using a Particle Flow jet composed of Particle Flow objects as described in section 2.3.3 are used.

In the Hadron plus Strips (HPS) algorithm, the number of electromagnetic objects is first estimated from the ECAL energy deposits. The algorithm then checks if the number of charged particles and electromagnetic objects and their momenta is compatible with one of the main tau decay modes. The tau visible momentum is then reconstructed as the sum of the various objects in the jet. Finally an isolation criteria is applied to reduce the contamination of fakes: the sum of transverse momenta of particles other than those connected to the decay and associated to the same primary vertex is required to be below a given value. The condition on the primary vertex allows one to discard charged particles coming from pile-up interactions. The neutral component from pile-up to the isolation sum is estimated in each event by multiplying the transverse momentum sum of charged particles not associated to the primary vertex, by a factor $\delta\beta$ representing the average ratio of neutral particles over charged particles in a pile-up interaction. This neutral component then subtracted to the final isolation sum.

Alternatively, a multivariate analysis is performed using variables from the Particle Flow objects. The efficiency for the HPS selection measured in the data and in simulation is shown in figure 2.19. Three working points are considered.

The tau reconstruction is hard in several aspects: the background from jet is important and, since at least one neutrino is produced during the tau decay, only a fraction of its total momentum is measured. The consequence is that in most analyses involving leptons in the final states, the electron and muon channel provides a much better sensitivity than the tau channel. In the search for heavy dilepton resonances for instance the use of the visible mass of the tau pair (defined as the invariant mass of the visible decay products of the tau with the missing transverse energy) leads to a width over mass ratio up to 50% for 2 TeV masses[75]. An important exception is the search for the scalar boson where

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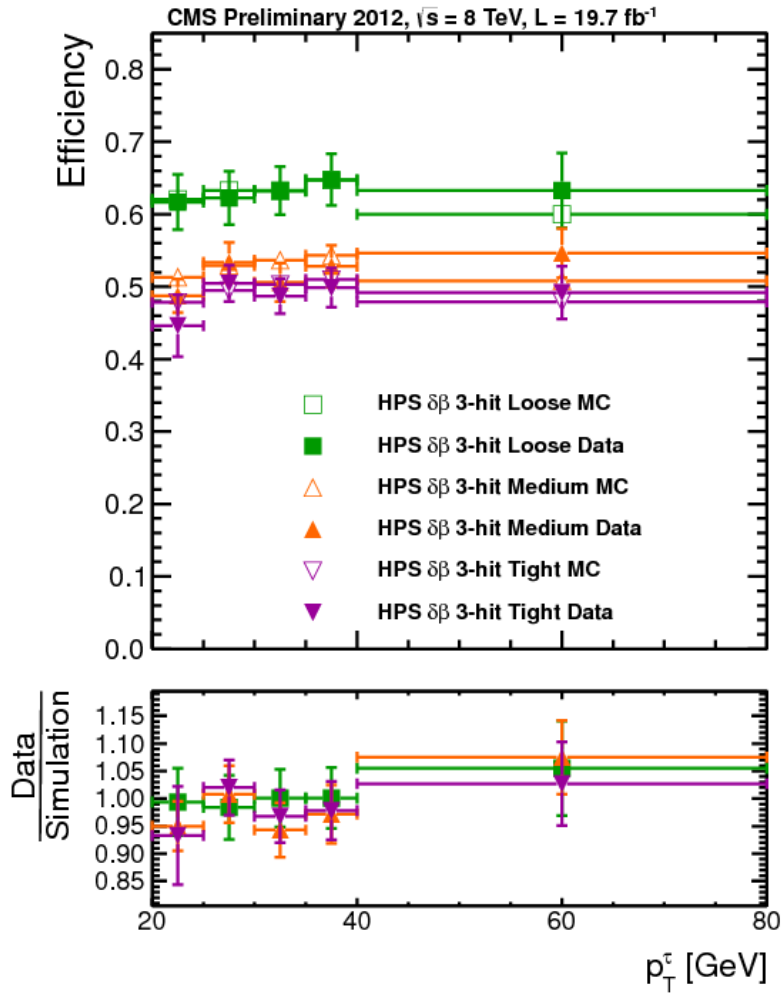


Figure 2.19: Efficiency of the Hadron plus Strips algorithm used for identifying a tau lepton decaying hadronically as a function of the tau lepton p_t . Results are presented for three working points for data and a simulated Drell-Yan sample. The bottom part shows the ratio of the data and simulation measurements. Figure is taken from [74].

the large enhancement of the cross section in this channel compared to the lighter leptons makes it from far the most sensitive channel.

2.3.6 Missing transverse energy

In a proton collision, the hard process (discussed in section 3.1.2), where particles with large transverse momentum are created, involves only one quark or gluon of each proton. The proton remnants will lead to the production of forward particles emitted mainly outside the detector acceptance. The consequence is that the total energy of the proton-proton collision cannot be measured. The transverse energy balance can on the other side be measured and the Particle Flow missing transverse energy (PFMet) is defined as the negative vectorial sum of the momentum of all the Particle Flow objects reconstructed in the event. Since the only known particles that escape detection are neutrinos, a large missing transverse energy is a hint for the production of such particles.

The measurement of the PFMet is affected by the experimental energy resolution on the various objects recorded in the event (especially jets). The PFMet distribution measured in part of the data recorded at $\sqrt{s} = 8$ TeV for events where two muons are reconstructed is shown in figure 2.20. This selection selects mainly Drell-Yan events where no neutrinos are expected and whose average PFMet value provides an estimate of the PFMet resolution, and leptonic decays of $t\bar{t}$ pairs where the two neutrinos coming from the top quark decay generates a large PFMet.

The PFMet is also of particular interest for the search of new particles with a small probability of interaction with ordinary matter (such as dark matter or supersymmetric particles).

2.4 8 TeV datasets

In this thesis the complete dataset recorded by the CMS detector for proton-proton collisions at $\sqrt{s} = 8$ TeV is analyzed. The recorded events are divided into 4 run periods (runs A, B, C, D) and sorted in various datasets depending on the trigger used to record them. The offline event reconstruction is processed several times. A prompt event reconstruction is first performed soon after the data are recorded and the data are then reprocessed once updated information about the detector (energy calibration, alignment,...) becomes available. The datasets used in most of the steps of the search for new heavy dielectron resonances (including the final mass spectrum measurement and the electron selection efficiency measurement) are shown in table 2.3. They correspond to the January 2013 reprocessing. The total integrated luminosity was estimated to be 19.7 fb^{-1} .

2. THE EXPERIMENTAL SETUP

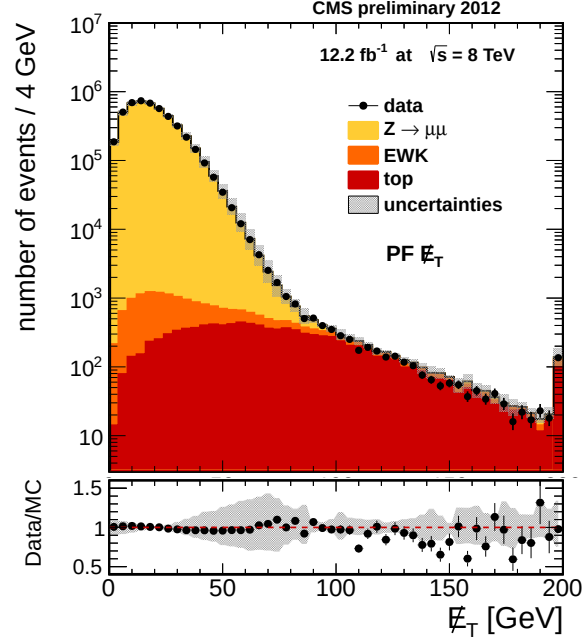


Figure 2.20: Distribution of the Particle Flow missing transverse energy in events with two reconstructed in data (black dots) and in various simulated samples (histograms) at $\sqrt{s} = 8$ TeV. Figure is taken from [76].

/Photon/Run2012A-22Jan2013-v1/AOD
/DoublePhotonHighPt/Run2012B-22Jan2013-v1/AOD
/DoublePhotonHighPt/Run2012C-22Jan2013-v1/AOD
/DoublePhotonHighPt/Run2012D-22Jan2013-v1/AOD

Table 2.3: The datasets used in the main parts of the search for heavy dielectron resonances at $\sqrt{s} = 8$ TeV.

Chapter 3

Simulation of proton-proton collisions

High energy proton-proton collisions provide an attractive framework to study Nature at the subatomic level and to understand its more fundamental laws. Extracting information from the data is however a challenging task that faces multiple issues, from both theoretical and experimental points of view.

The first difficulty arises from the nature of the protons themselves. It is known since more than half a century that they are not fundamental particles and possess a substructure described as a complicated entanglement of quarks and gluons generically called partons. At high energy, proton-proton collisions should therefore rather be described in term of quarks and gluons collisions.

Since the subatomic level is governed by the laws of quantum physics, the outcome of a proton-proton interaction can never be predicted exactly and a probabilistic description which takes into account all the possible outcomes must instead be adopted. The vast majority of proton-proton interactions consists of elastic scattering or inelastic scattering with a limited momentum transfer (soft scattering). In particle physics, the more interesting collisions are however those involving a large exchange of momentum (hard scattering) which are much rarer. Figure 3.1 illustrates this fact and shows the total proton-proton or proton-antiproton cross section and the cross sections estimated for various processes such as the production of a W or a Z boson, as a function of the center of mass energy, $\sqrt{s}$. At $\sqrt{s} = 8$ TeV, the total proton-proton production cross section has been measured by the TOTEM collaboration to be 101.7 ± 2.9 mb [78]. The W or Z boson cross section is a factor 10^6 times smaller while for a scalar boson with a mass of 125 GeV the ratio is 10^{10} . In this thesis, processes with cross sections down to the femtobarn scale, that is, 10^{14} times smaller than the inclusive cross section, will be searched for. Even between the different hard processes, a large hierarchy exists, and in a hadron collider the production of hard jets represents the first contribution to the hard scattering cross section by several orders of magnitude.

The hard inelastic scattering is modelled in two parts: the hard interaction between two partons, where perturbative quantum theory calculations (*cf.* section 1.1.4) can be used, and the recombination of the protons remnants. Perturbative calculations such as the one detailed in appendix A predict the hard interaction kinematics with great

3. SIMULATION OF PROTON-PROTON COLLISIONS

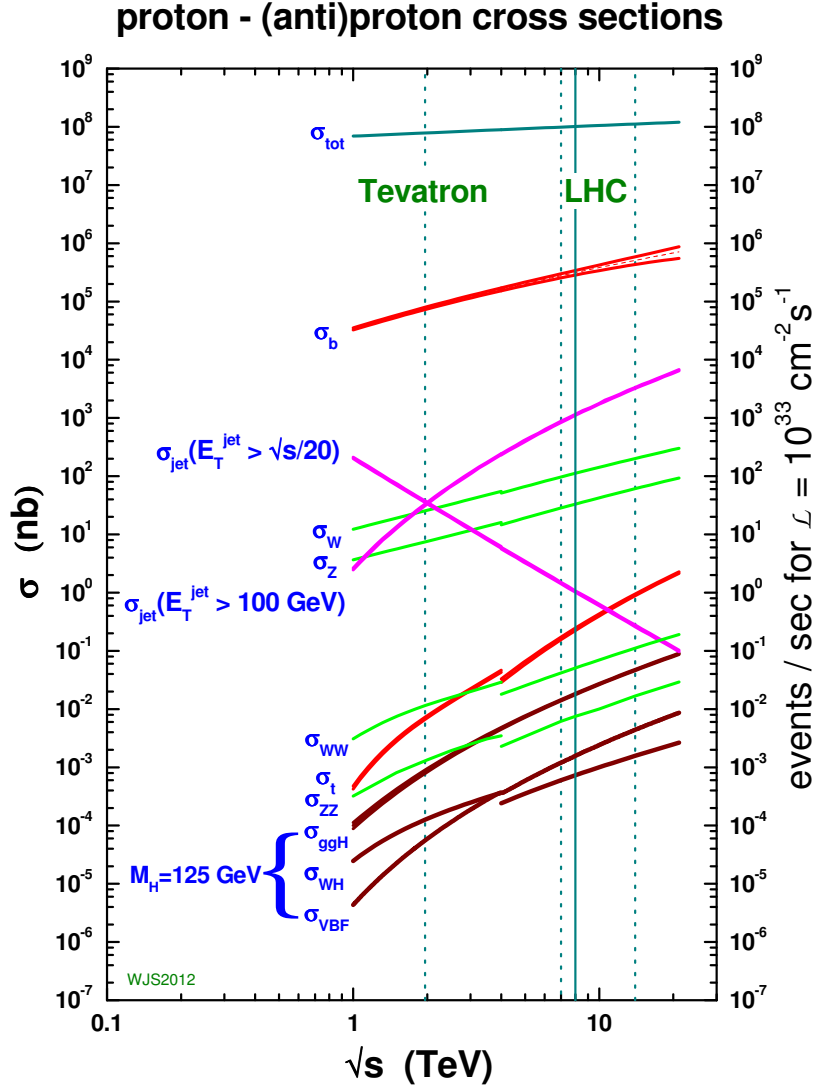


Figure 3.1: The total proton-proton or proton-antiproton cross section and the cross sections estimated for various processes (production of bottom or top quarks, a Z or W boson, a scalar boson, or jets above a transverse momentum threshold) as a function of the center of mass energy, $\sqrt{s}$. Figure is taken from [77].

accuracy. However, the treatment of real and virtual radiative corrections is complex and various approaches with their specific pros and cons are currently used. Since the hard interaction deals with quarks and gluons which are confined in hadrons, a model describing their hadronisation is also needed.

The sophistication of the experimental apparatus implies that, to the complexity of the theoretical description of a collision, must be added the experimental complexity of its reconstruction by the detector.

To overcome those difficulties, the proton-proton collisions description relies heavily on Monte Carlo simulations: a large number of individual events are first generated and the interaction of the various particles present in the final state with the detector is then simulated.

This chapter explains how such collisions are generated and how the detector response is simulated.

3.1 Generation of proton-proton collisions

The generation of hard inelastic proton-proton collisions follows several steps depicted in figure 3.2 and described in the following subsections.

- To estimate the cross section associated to a hard process involving a parton A from the first proton and a parton B from the second proton, one first needs to know the probability of finding such partons in each proton and their momentum distribution. This is achieved with the help of the parton distribution functions defined in section 3.1.1.
- The cross section is then factorized as the product of the parton distribution functions and the parton-parton cross section, calculated from the scattering matrix element. The matrix element calculation, which can involve the leading order contribution or in some cases next to leading order contribution is discussed in 3.1.2.
- Quarks and gluons in the initial and final states are subject to radiation. This effect is accounted for by using a so-called Parton Showering technique described in section 3.1.3.
- The interaction between the rest of the two protons forms the underlying event and is the topic of section 3.1.4.
- Below some energy scale Λ_{QCD} where the strong coupling constant α_s becomes comparable or greater than 1, quarks and gluons, which carry colour, cannot be considered as free anymore and recombine to form neutral hadrons. This hadronisation phenomenon is described in section 3.1.5.
- At the LHC, several collisions happen during the same bunch crossing and the hard collision of interest is superimposed to other softer collisions. The simulation of these pile-up events is explained in section 3.1.6.

3. SIMULATION OF PROTON-PROTON COLLISIONS

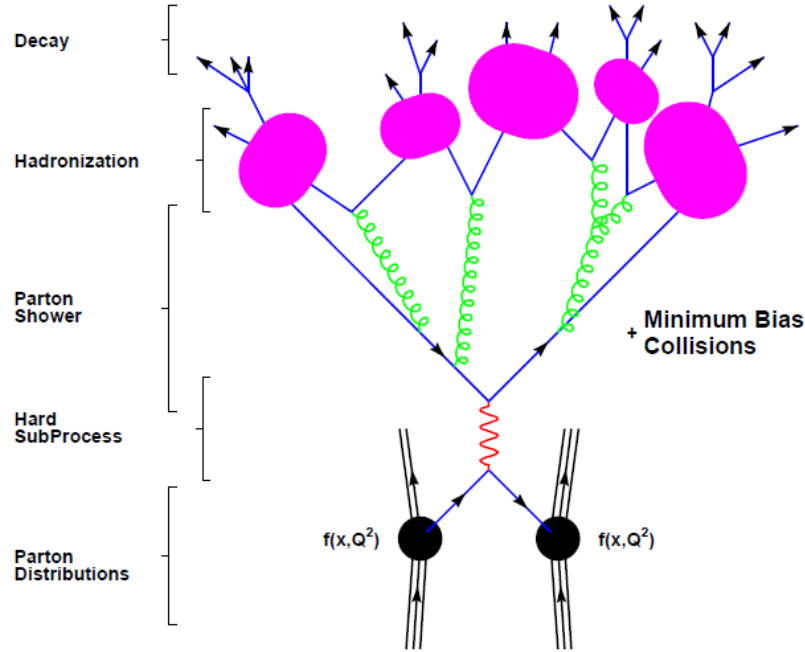


Figure 3.2: The various steps involved in the generation of a proton-proton collision.

3.1.1 The parton distribution functions

In proton-proton collisions involving a large 4-momentum transfer squared Q^2 , the parton distribution functions (pdf) represent the density of a given parton with a fraction x of the total proton momentum in the proton-proton center of mass frame. They are not predicted theoretically and must be experimentally measured. Currently, they are essentially extracted from the study of lepton-hadron collisions such as those provided by HERA [79, 80]. At hadron colliders, some constraints are also obtained by the study of the Drell-Yan process [81] or by the measurement of the W^+/W^- asymmetry [82]. An important feature of the pdf is that they depend on the energy scale characterizing the hard process: due to quantum fluctuations, the number of visible gluons and quarks increases with the energy scale at which the proton is probed quantified by the quantity Q^2 . The low energy picture of the proton, described as composed of three valence quarks, does not hold for high Q^2 and one must add to the valence, the presence of a sea of quarks and gluons.

The pdf measured at a given Q^2 can be extrapolated at another energy using the Dokshitzer-Gribov-Lipatov-Altarelli-Parisi (DGLAP) equation [83, 84, 85]. One first de-

defines the splitting functions P_{ba} :

$$\begin{aligned}
P_{qq}(z) &= \frac{4}{3} \left(\frac{1+z^2}{1-z} + \frac{3}{2} \delta(1-z) \right) \\
P_{qg}(z) &= \frac{1}{2} (z^2 + (1-z)^2) \\
P_{gq}(z) &= \frac{4}{3} \left(\frac{1+(1-z)^2}{z} \right) \\
P_{gg}(z) &= 6 \left(\frac{z}{1-z} + \frac{1-z}{z} + z(1-z) + \frac{33-2n_f}{6} \delta(1-z) \right)
\end{aligned} \tag{3.1}$$

where n_f is the number of quark flavours relevant at the scale considered. The P_{ba} are related to the probability for a parton a to branch into two partons b and c , where parton b carries a fraction z of its momentum. The following processes are possible: $q \rightarrow qq$ (P_{qq}), $g \rightarrow qq$ (P_{gq}) and $g \rightarrow gg$ (P_{gg})¹. The probability for a parton a with a virtuality Q^2 to branch into two parton b and c is given by the DGLAP equation:

$$d\mathcal{P}_{a \rightarrow bc} = \frac{dQ^2}{Q^2} \frac{\alpha_s}{2\pi} dz P_{ba}(z) \tag{3.2}$$

which can be used to derive the evolution of the pdf (at first order in the strong coupling constant α_s):

$$Q^2 \partial_{Q^2} \begin{pmatrix} q_i(x, Q^2) \\ \bar{q}_i(x, Q^2) \\ g(x, Q^2) \end{pmatrix} = \frac{\alpha_s(Q^2)}{2\pi} \sum_{j=1}^{n_f} \int_x^1 \frac{d\xi}{\xi} \begin{pmatrix} P_{qq}(\frac{x}{\xi}) \delta_{ij} & 0 & P_{qg}(\frac{x}{\xi}) \\ 0 & P_{qq}(\frac{x}{\xi}) & P_{qg}(\frac{x}{\xi}) \\ P_{gq}(\frac{x}{\xi}) & P_{gq}(\frac{x}{\xi}) & P_{gg}(\frac{x}{\xi}) \end{pmatrix} \begin{pmatrix} q_i(x, Q^2) \\ \bar{q}_i(x, Q^2) \\ g(x, Q^2) \end{pmatrix} \tag{3.3}$$

where $q_i(x, Q^2)$, $\bar{q}_i(x, Q^2)$ and $g(x, Q^2)$ are the pdf associated to the quark or antiquark of flavour i and to the gluon respectively. Those equations allow one to extrapolate the pdf from a scale μ_0^2 where they were measured to an arbitrary scale μ_F^2 called the factorization scale. This scale fixes the maximal radiation virtuality treated by the evolution of the pdf. Higher virtualities must on the other side be included in the matrix element calculation.

Various sets of pdf exist and are provided by different collaborations such as CTEQ [86, 87], MSTW [88] and NNLOPDF [89]. The pdf parametrization in the CT10 pdf sets released by the CTEQ group is shown in figure 3.3 for several quarks and for the gluon. In the figure, the pdf are multiplied by the fraction of the total momentum carried by the parton x for normalization purposes. The bump around $x = 0.1$ visible for up and down quarks in the left plot of figure 3.3 corresponds to the contribution of valence quarks. For x values below 0.01 however, the sea quarks dominate. In this thesis, a search for TeV resonances is performed. The Q^2 values considered are therefore the invariant mass squared of the system, M^2 , which can be related to the fractions of momentum carried by the two colliding partons x_1 and x_2 through the relation:

$$M = \sqrt{x_1 x_2 s} \tag{3.4}$$

¹quarks and antiquarks are both denoted q here.

3. SIMULATION OF PROTON-PROTON COLLISIONS

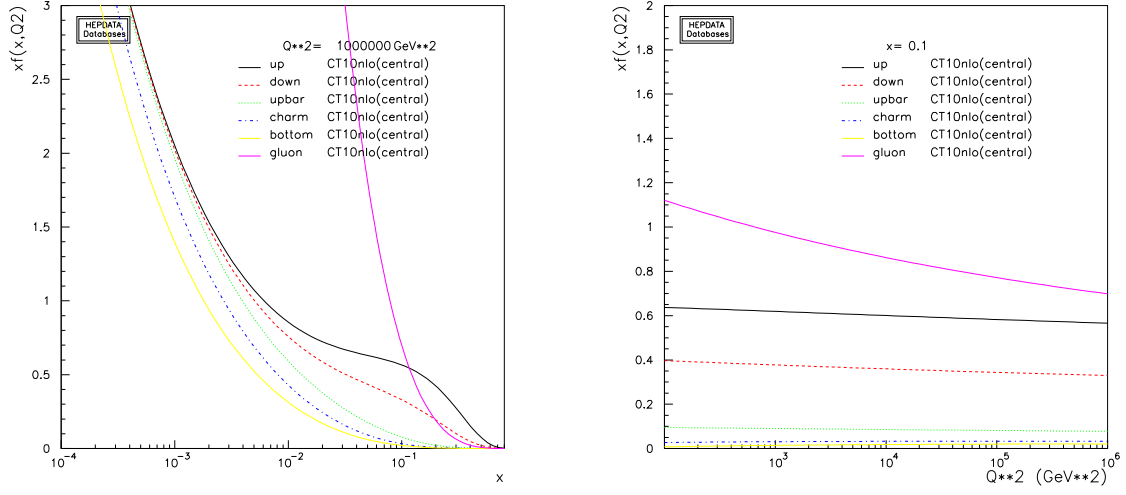


Figure 3.3: The parton distribution function (pdf) times the fraction of the proton energy carried by the quark, x as given by the CT10 parametrization. The pdf are shown for up, down, charm and bottom quarks, up antiquark and gluon and are presented as a function of x for $Q^2 = 1 \text{ TeV}^2$ (left) and as a function of Q^2 for $x = 0.1$ (right). The plots were generated using the pdf plotter provided in [90].

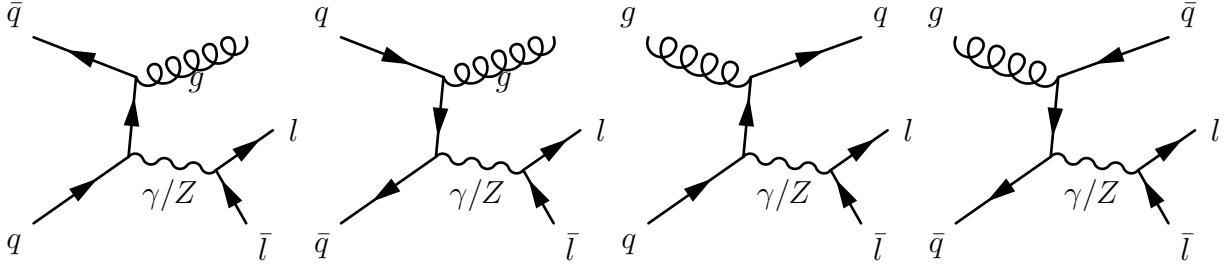


Figure 3.4: Feynman diagrams corresponding to the first order real radiative corrections to the Drell-Yan process.

For a mass in the TeV range and $\sqrt{s} = 8$ TeV, the resulting average x is around 0.1. At leading order, the two initial partons involved in the production of a neutral resonance are a quark-antiquark or a gluon-gluon pair. In the first case, due to the valence contribution, the quark carries in average a larger x than the antiquark. The consequence is that a strong boost of the partonic system is expected along the z axis in the laboratory frame. For the gluon-gluon case, the collision is more symmetric.

A property of the sea quarks and gluon pdf is that for a given change of $x \rightarrow a.x$ with $a > 1$, the ratio $f(x, Q^2)/f(a.x, Q^2)$ increases with x . Equation 3.4 therefore implies that the production cross section of a TeV resonance at the LHC is extremely sensitive to $\sqrt{s}$.

The parametric functions used to model the pdf are only known with a certain precision and their extrapolation from Q^2 values where they were measured to the region of interest is therefore subject to possibly large uncertainties. The uncertainty on the cross section of a given process due to the pdf is therefore estimated by consistently varying all the pdf parameters by one standard deviation.

3.1.2 Hard process

The hard process part of the proton-proton collision involves only two partons a and b whose interaction is characterized by a large energy scale Q^2 . The strong coupling constant is well below 1 and a perturbative approach can therefore be adopted. It consists in computing the matrix element between the initial state and a final state X as the sum of contributions with increasing powers of α_s to finally derive the cross section associated to the process, $\sigma_{ab \rightarrow X}$.

For instance, the leading contribution to the Drell-Yan process can be calculated from the diagrams already introduced in figure 1.2. An accurate description of the process must however also take into account radiative corrections. They can be of two types: real (production of a quark or a gluon in addition to the photon or the Z boson) or virtual (loop correction to the vertex or to the quark or antiquark propagator). The diagrams corresponding to the real and virtual corrections at the first order are respectively shown in figure 3.4 and figure 3.5.

Once the $ab \rightarrow X$ partonic cross section is known, the total proton-proton cross section can be factorized [91] as:

$$\sigma_{pp \rightarrow X} = \sum_{a,b=\{q,\bar{q},g\}} \int_0^1 dx_1 dx_2 f_a(x_1, \mu_F) f_b(x_2, \mu_F) \sigma_{ab \rightarrow X}(x_1, x_2, \mu_F, \alpha_s(\mu_R)) \quad (3.5)$$

3. SIMULATION OF PROTON-PROTON COLLISIONS

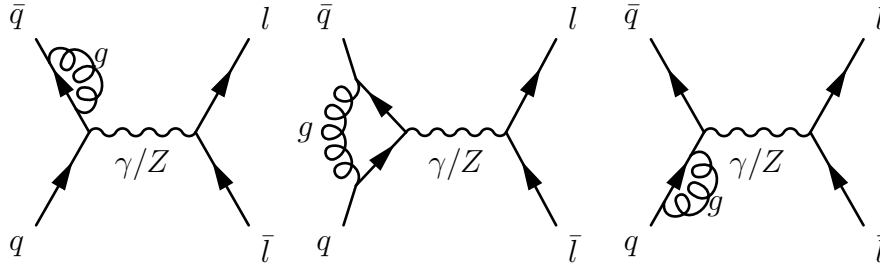


Figure 3.5: Feynman diagrams corresponding to the first order virtual radiative corrections to the Drell-Yan process.

where α_R is the renormalized value of the strong coupling constant at a chosen scale μ_R .

Various generators are available to generate events according to a given hard process [92]. Although in principle all the orders of the perturbative development should be considered, most of them only include the leading contribution to the matrix element.

A popular generator is Pythia [93, 94], a general purpose program which, in addition to the hard process, also takes care of the parton showering, the hadronisation, and the description of the underlying event. For the matrix element calculation, Pythia only considers the leading order (diagrams in figure 1.2 in the Drell-Yan case). The radiative corrections are not explicitly calculated but rather treated in the parton showering. As will be discussed in the next section, the parton showering is well suited to treat low Q^2 radiation. It however fails to accurately describe hard radiations. The consequence is that the performances of Pythia to describe events including many hard jets in the final state are limited.

A more accurate approach is followed by MadGraph [95] where the hard real radiative corrections are included in the matrix element (figure 3.5). Diagrams with up to 5 extra partons are included. This generator is well suited to study processes such as W or Z produced in addition with hard jets.

Finally, the Powheg [96] generator is a complete next to leading order (NLO) generator that considers both real and virtual corrections at the leading order. It is therefore limited to the radiation of one hard extra parton after what it must rely on the parton showering description. The use of the full NLO nevertheless provides in general an accurate estimate of the inclusive cross section. Since the first parton radiation is already described, the interface between Powheg and the parton showering requires some special treatment.

In addition to event generators, there also exist cross section integrators which can be used to compute cross sections as a function of a variable of interest. An example is the FEWZ code [97] which allows one to make calculation at the NNLO level for exclusive production of a W or Z boson.

3.1.3 Parton showering

Partons involved in the initial and final states of the hard process are subject to radiation. Their incorporation in the matrix element is possible, and even desirable for hard radiation, but the current generators are only able to treat a limited number of partons in the final state. Moreover, the matrix element calculation leads to divergencies for collinear or soft radiation if one does not consider both real and virtual contributions.

A Parton Showering technique is therefore adopted to describe the parton radiation from a virtuality $Q^2 = \mu_F^2$ down to a cut off scale $Q_{cutoff}^2 \approx 1 \text{ GeV}^2$ at which hadronisation occurs. Equation 3.2 implies that the probability $\mathcal{P}_{a \rightarrow a}$ for a parton a to evolve without radiating from a virtuality Q_0^2 to a virtuality Q^2 is given by the Sudakov form factor [98]:

$$\mathcal{P}_{a \rightarrow a}(Q^2, Q_0^2) = \exp \left(- \sum_b \int_{Q^2}^{Q_0^2} \frac{dQ^2}{Q^2} \frac{\alpha_s}{2\pi} \int P_{ab}(z) dz \right) \quad (3.6)$$

For a given event, a random number R is drawn from an uniform distribution between 0 and 1 and a first branching is simulated at a Q^2 value such that $\mathcal{P}_{a \rightarrow a}(Q^2, Q_0^2) = R$ if $Q^2 > Q_{cutoff}^2$. The showering is then repeated for the daughter partons until Q_{cutoff}^2 is reached.

The showering procedure is applied to partons both in the initial and final states, the main difference being the sign of the virtualities considered (negative for the initial state since one goes back in time).

3.1.4 Underlying event

In addition to the hard process, the interactions between the rest of the protons must also be considered. Since the parton participating to the hard process is missing, the proton remnants are not colorless and will hadronize. The produced hadrons will however carry a very small transverse momentum and will be very forward.

It is also possible that more than one parton interact with the other proton. This phenomenon, called multiple parton interaction, has already been observed when one of the interaction is soft but its description is still a field of research [99]. The instantaneous luminosities reached at the LHC should allow one to observe multiple interactions involving two hard processes which should shed light on the proton structure.

The description of the underlying event uses phenomenological models that are tuned on the data via the measurement of the multiplicity of low transverse momentum tracks [100].

3.1.5 Hadronisation

For energies approaching Λ_{QCD} , where the strong coupling constant becomes close to 1, the perturbative approach does not hold anymore and the various partons produced in the previous steps recombine to form colorless states.

This phenomenon is currently described only phenomenologically through various models. A popular one, called the Lund string model [101] and used by Pythia, describes partons moving away from each other as connected by a string with a fixed energy per unit length of the order 1 GeV/fm. When this energy becomes large enough, the string breaks down and creates a quark-antiquark pair, cancelling out the color field between the initial partons which form hadrons. The probability for a parton to give rise to a specific hadron with a given fraction of the string momentum is modelled by empirical functions called fragmentation functions, whose exact expressions are not known from the theory. They are therefore adjusted to fit to the experimental measurements. The fragmentation process continues with the string remnants until all the partons have hadronized.

3. SIMULATION OF PROTON-PROTON COLLISIONS

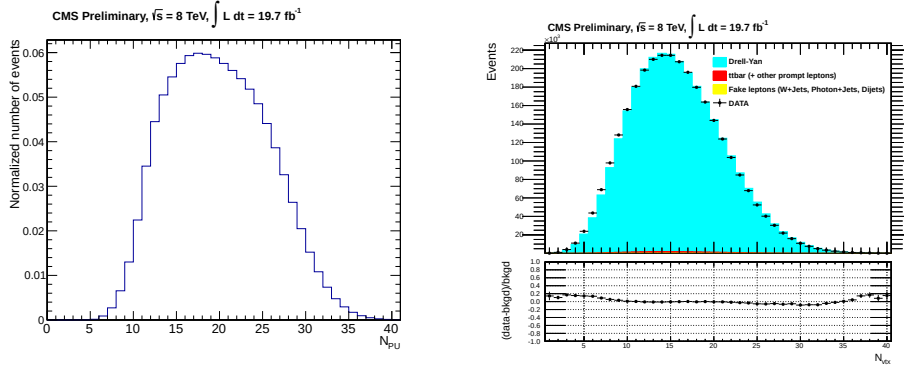


Figure 3.6: The distribution of the estimated average number of pile-up collisions in the full dataset of proton-proton collisions recorded at $\sqrt{s} = 8$ TeV in 2012 (left) and the distribution of the number of reconstructed primary vertices in data (black dots) and in various simulated samples after pile-up reweighting, for a dielectron sample enriched in Drell-Yan events (right).

3.1.6 Pile-up

In order to accumulate enough data to study rare processes, the collision rate at the LHC is high and for the data taken in 2012, an average of 20 proton-proton inelastic interactions occurred per bunch crossing. The consequence is that particles coming from different interactions are superimposed in the detector. This effect is taken into account in the simulations and in addition to the collision producing the hard process under study, the products of several soft inelastic collisions are also added.

The exact number of average collisions per bunch crossing in the data can be estimated by multiplying the instantaneous luminosity, continuously monitored, by the total inelastic cross section. One can then reconstruct the distribution of the number of pile-up interactions in the data for the complete data taking. The corresponding distribution for the 2012 data is shown on left part of figure 3.6. Simulated events are then reweighted such that they match the data distribution.

The description of the pile-up by the simulation can be verified by counting the number of reconstructed vertices in the event as illustrated in the right part of figure 3.6.

3.2 Detector simulation

After a collision has been generated, one still must simulate the detector response. For the CMS detector, it is the GEANT4 [102] program which calculates the trajectory of the various particles generated during the collision, simulates their interaction with matter and the signal they will induce in the various subdetectors. Amongst other things, GEANT4 is in charge of describing the bremsstrahlung emissions due to interaction of electrons with the tracker layers. The detector geometry is given as an input to the program and to obtain a description as close as possible to the reality, any available information such

as the existence of insensitive materials or dead channels and their position, is included.

After this step, a simulated collision has exactly the same format than a collision recorded in the data. The same code is then used to reconstruct the various physical objects described in section 2.3.

The software used by CMS, CMSSW, is a code mainly written in the C++ language which is constantly updated by the whole collaboration to integrate new knowledge on the detector (new tracker alignment or calorimeter calibration,...) or to improve the object reconstruction.

3.3 Simulated samples and datasets

Tables 3.1 and 3.2 present the list of the various background and signal simulated samples that were used in this thesis. Table 3.1 includes a quite exhaustive list of the Standard Model processes. The signal samples in table 3.2 were generated without taking the interference with the Drell-Yan process into account. They are namely the Z'_{SSM} and Z'_Ψ (*cf.* section 1.4.2.2) and the Randall-Sundrum graviton G_{RS} (*cf.* section 1.4.3.3). All the samples were reconstructed with the version 53X of the CMSSW software.

The process cross section, the event generator and pdf sets used for the generation are also indicated in the table. The fragmentation and hadronisation steps are described using Pythia in all the samples. In most of the cases, the cross section obtained by the generator is replaced by a cross section calculation taking into account higher order corrections. This is especially the case for the Drell-Yan process and the W production which are rescaled to the NNLO cross section measured by the FEWZ calculator [97] at the Z peak. The $t\bar{t}$ and tW cross sections can be found in references [103] and [104] respectively. Finally, the diboson and γ +jets cross sections were calculated with the MCFM [105] and DIPHOX [106] calculator respectively.

When mentioned, the p_t range indicates the range allowed for the outgoing particles in the matrix element calculation.

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Sample	Cross-section (pb)	Generator	Pdf set
Drell-Yan ($e\bar{e}$ channel) $M_{e\bar{e}} > 20$	1915 (NNLO)	Powheg	CT10
Drell-Yan ($e\bar{e}$ channel) $M_{e\bar{e}} > 120$	12.16 (NLO $\times$ 1.024)	Powheg	CT10
Drell-Yan ($e\bar{e}$ channel) $M_{e\bar{e}} > 200$	1.517 (NLO $\times$ 1.024)	Powheg	CT10
Drell-Yan ($e\bar{e}$ channel) $M_{e\bar{e}} > 400$	0.1112 (NLO $\times$ 1.024)	Powheg	CT10
Drell-Yan ($e\bar{e}$ channel) $M_{e\bar{e}} > 500$	0.04515 (NLO $\times$ 1.024)	Powheg	CT10
Drell-Yan ($e\bar{e}$ channel) $M_{e\bar{e}} > 700$	0.01048 (NLO $\times$ 1.024)	Powheg	CT10
Drell-Yan ($e\bar{e}$ channel) $M_{e\bar{e}} > 800$	0.005615 (NLO $\times$ 1.024)	Powheg	CT10
Drell-Yan ($e\bar{e}$ channel) $M_{e\bar{e}} > 1000$	0.001837 (NLO $\times$ 1.024)	Powheg	CT10
Drell-Yan ($e\bar{e}$ channel) $M_{e\bar{e}} > 1500$	$1.744 \cdot 10^{-4}$ (NLO $\times$ 1.024)	Powheg	CT10
Drell-Yan ($e\bar{e}$ channel) $M_{e\bar{e}} > 2000$	$2.259 \cdot 10^{-5}$ (NLO $\times$ 1.024)	Powheg	CT10
Drell-Yan ($\mu\bar{\mu}$ channel) $M_{\mu\bar{\mu}} > 20$	1915 (NNLO)	Powheg	CT10
Drell-Yan ($\tau\bar{\tau}$ channel) $M_{\tau\bar{\tau}} > 20$	1915 (NNLO)	Powheg	CT10
$t\bar{t}$	245.8 (NNLO)	Powheg	CT10
$t\bar{t}$, $700 < M_{t\bar{t}} < 1000$	18.2 (NLO $\times$ 1.17)	Powheg	CT10
$t\bar{t}$, $M_{t\bar{t}} > 1000$	3.4 (NLO $\times$ 1.17)	Powheg	CT10
tW	11.1 (NNLO)	Powheg	CT10
$\bar{t}W$	11.1 (NNLO)	Powheg	CT10
WW	54.8 (NLO)	Pythia 6	CTEQ6
WZ	33.2 (NLO)	Pythia 6	CTEQ6
ZZ	17.7 (NLO)	Pythia 6	CTEQ6
$W \rightarrow l\nu$	36257 (NNLO)	Madgraph	CTEQ6
γ +jets, $15 < p_t < 30$	200062×1.3 (LO $\times$ 1.3)	Pythia 6	CTEQ6
γ +jets, $30 < p_t < 50$	19932×1.3 (LO $\times$ 1.3)	Pythia 6	CTEQ6
γ +jets, $50 < p_t < 80$	3322.3×1.3 (LO $\times$ 1.3)	Pythia 6	CTEQ6
γ +jets, $80 < p_t < 120$	558.3×1.3 (LO $\times$ 1.3)	Pythia 6	CTEQ6
γ +jets, $120 < p_t < 170$	108.0×1.3 (LO $\times$ 1.3)	Pythia 6	CTEQ6
γ +jets, $170 < p_t < 300$	30.12×1.3 (LO $\times$ 1.3)	Pythia 6	CTEQ6
γ +jets, $300 < p_t < 470$	2.139×1.3 (LO $\times$ 1.3)	Pythia 6	CTEQ6
γ +jets, $470 < p_t < 800$	0.2119×1.3 (LO $\times$ 1.3)	Pythia 6	CTEQ6
γ +jets, $800 < p_t < 1400$	0.007078×1.3 (LO $\times$ 1.3)	Pythia 6	CTEQ6
γ +jets, $1400 < p_t < 1800$	$4.510 \cdot 10^{-5} \times 1.3$ (LO $\times$ 1.3)	Pythia 6	CTEQ6
γ +jets, $p_t > 1800$	$1.867 \cdot 10^{-6} \times 1.3$ (LO $\times$ 1.3)	Pythia 6	CTEQ6
QCD double $30 < p_t < 80$	$7.433 \cdot 10^7$ (LO)	Pythia 6	CTEQ6
QCD double $80 < p_t < 170$	$1.191 \cdot 10^6$ (LO)	Pythia 6	CTEQ6
QCD double $170 < p_t < 250$	30990 (LO)	Pythia 6	CTEQ6
QCD double $250 < p_t < 350$	4250 (LO)	Pythia 6	CTEQ6
QCD double $p_t > 350$	810 (LO)	Pythia 6	CTEQ6

Table 3.1: Background Monte Carlo samples used in this analysis, together with the cross-sections used for the process. All the masses and p_t are indicated in GeV and GeV respectively.

Sample	Cross-section (pb)	Generator	Pdf set
$Z'_\Psi \rightarrow e\bar{e}$ M = 750	0.1328 (LO)	Pythia 6	CTEQ6
$Z'_\Psi \rightarrow e\bar{e}$ M = 1000	0.03933 (LO)	Pythia 6	CTEQ6
$Z'_\Psi \rightarrow e\bar{e}$ M = 1250	0.01196 (LO)	Pythia 6	CTEQ6
$Z'_\Psi \rightarrow e\bar{e}$ M = 1500	0.00437 (LO)	Pythia 6	CTEQ6
$Z'_\Psi \rightarrow e\bar{e}$ M = 1750	0.00168 (LO)	Pythia 6	CTEQ6
$Z'_\Psi \rightarrow e\bar{e}$ M = 2000	$7.029 \cdot 10^{-4}$ (LO)	Pythia 6	CTEQ6
$Z'_\Psi \rightarrow e\bar{e}$ M = 2250	$2.895 \cdot 10^{-4}$ (LO)	Pythia 6	CTEQ6
$Z'_\Psi \rightarrow e\bar{e}$ M = 3000	$2.666 \cdot 10^{-5}$ (LO)	Pythia 6	CTEQ6
$Z'_{SSM} \rightarrow e\bar{e}$ M = 2250	0.001302 (LO)	Pythia 6	CTEQ6
$G_{RS} \rightarrow e\bar{e}$ M = 250	1.11 (LO)	Pythia 6	CTEQ6
$G_{RS} \rightarrow e\bar{e}$ M = 500	0.04303 (LO)	Pythia 6	CTEQ6
$G_{RS} \rightarrow e\bar{e}$ M = 750	0.005059 (LO)	Pythia 6	CTEQ6
$G_{RS} \rightarrow e\bar{e}$ M = 1000	$9.653 \cdot 10^{-4}$ (LO)	Pythia 6	CTEQ6
$G_{RS} \rightarrow e\bar{e}$ M = 1250	$2.478 \cdot 10^{-4}$ (LO)	Pythia 6	CTEQ6
$G_{RS} \rightarrow e\bar{e}$ M = 1500	$7.303 \cdot 10^{-5}$ (LO)	Pythia 6	CTEQ6
$G_{RS} \rightarrow e\bar{e}$ M = 1750	$2.584 \cdot 10^{-5}$ (LO)	Pythia 6	CTEQ6
$G_{RS} \rightarrow e\bar{e}$ M = 2000	$9.288 \cdot 10^{-6}$ (LO)	Pythia 6	CTEQ6
$G_{RS} \rightarrow e\bar{e}$ M = 2250	$3.458 \cdot 10^{-6}$ (LO)	Pythia 6	CTEQ6
$G_{RS} \rightarrow e\bar{e}$ M = 3000	$2.368 \cdot 10^{-7}$ (LO)	Pythia 6	CTEQ6

Table 3.2: Signal Monte Carlo samples used in this analysis, together with the cross-sections used for the process. All the masses and p_t are indicated in GeV and GeV respectively.

3. SIMULATION OF PROTON-PROTON COLLISIONS

Chapter 4

Reconstruction and selection of high energy electrons with the CMS detector

This chapter is devoted to the reconstruction and selection of high energy electrons with the CMS detector. The first two sections (4.1, 4.2) describe the general reconstruction of electrons candidates (denoted GSF electron candidates) with the detector, and how they are triggered. Section 4.3 constitutes the core of this chapter and details a set of selection criteria optimized for high energy electrons. It is shown in section 4.4 that some of those criteria are affected by the presence of pile-up collisions and a correction for this effect is presented. Section 4.5 summarizes the results presented in this chapter and presents the final selection efficiency for electrons and fake candidates from jets.

4.1 Electron reconstruction

As briefly described in section 2.3.2, the electron reconstruction is based on the association of an energy deposit in the ECAL with a track reconstructed in the tracker detector. Given the tracker acceptance in CMS (*cf.* section 2.2.4), the electron reconstruction is therefore possible for pseudorapidity in the range $(-2.5, 2.5)$. One of the specificities of the electrons is their high probability of emitting a bremsstrahlung photon after interacting with one of the tracker layers. This feature is illustrated in figure 4.1 that shows, for electrons with transverse momentum, p_t , of 20 GeV or more, the fraction of energy carried by bremsstrahlung photons, f_{brem} , defined as:

$$f_{brem} = \frac{p_{in} - p_{out}}{p_{in}} \quad (4.1)$$

where p_{in} and p_{out} are the momenta measured from the electron track at the innermost and outermost layers of the tracker respectively. The fact that the bremsstrahlung photons can carry a large fraction of the electron energy has two consequences. First, because of the bending of the electron trajectory in the magnetic field, the energy deposits of the bremsstrahlung photons and of the electron are spread in ϕ . Secondly, a kink in the

4. RECONSTRUCTION AND SELECTION OF HIGH ENERGY ELECTRONS WITH THE CMS DETECTOR

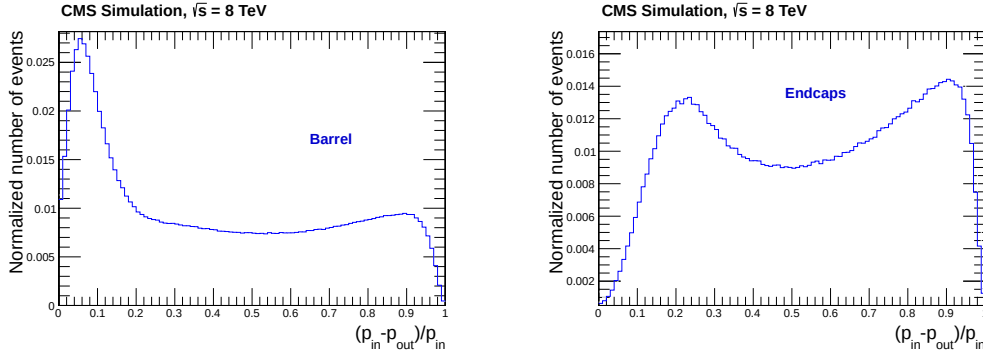


Figure 4.1: Distribution of the fraction of the electron momentum carried by photons radiated in the tracker, obtained by comparing the electron track momentum as measured in the innermost and outermost layers of the tracker, for electrons with $p_t > 20$ GeV coming from a simulated sample of Drell-Yan events with $M_{ee} > 20$ GeV (first line of table 3.1) emitted in the ECAL barrel (left) or endcaps (right).

reconstructed trajectory is expected in the layer where a photon is radiated. Those two effects need to be properly treated in the reconstruction algorithms [9].

4.1.1 ECAL clustering

The lateral shower shape of an electron entering the electromagnetic calorimeter (ECAL) is well known: in average, 94% (resp. 97%) of its energy is contained in a 3×3 (resp. 5×5) matrix of crystals. Recovering the energy carried by bremsstrahlung photons implies however to look for secondary deposits in a large ϕ window. Because of the two different geometries in the central and forward parts of the ECAL, this is achieved with two different algorithms [9, 107].

In the barrel, the so-called “hybrid” algorithm is used. It takes advantage of the fact that all crystals have a constant η and ϕ width: all crystals with a transverse energy deposit E_T above 1 GeV are used as seed, starting from the highest E_T one. A band of 3 crystals in η and 35 crystals in ϕ centered on the seed is constructed and divided into 35 dominos of constant ϕ . If a domino has an energy E above 1 GeV, it is extended 1 crystal up and down in η . A cluster is then made by assembling the seed domino with all dominos satisfying $E > 0.1$ GeV and that are adjacent to the seed domino or to a domino already added to the cluster. Secondary clusters are formed in the same way starting for any unused domino with $E > 0.35$ GeV. Finally, all the clusters in the band are gathered to form a supercluster (SC).

In the endcaps, the “multi 5×5 ” algorithm also starts from a seed crystal with $E_T > 0.18$ GeV, around which a 5×5 matrix is built, which constitutes the primary cluster. Secondary 5×5 matrices are then made centered on all neighbour crystals with $|\eta_{neighbour} - \eta_{seed}| < 0.07$ and $|\phi_{neighbour} - \phi_{seed}| < 0.3$. The supercluster is then constituted of the primary cluster and all the secondary matrices with a transverse energy above 1 GeV.

Finally, the energy of the preshower cells in front of the crystals included in the ECAL clusters is added to the supercluster.

The energy leakage outside the supercluster and in other subdetectors and the contribution from pile-up noise affect the energy measurement. Corrections to the supercluster energy are therefore derived using a multivariate technique trained on a simulation sample that takes as input variables such as the supercluster position, the HCAL over ECAL energy ratio and the average energy density in the whole detector.

The supercluster position $\vec{x}$ is obtained from a weighted sum of the individual crystals position $\vec{x}_i$ and energy E_i :

$$\vec{x} = \frac{\sum_i \vec{x}_i \omega_i}{\sum_i \omega_i} \quad (4.2)$$

where the weights ω_i are:

$$\omega_i = \max \left(\omega_0 + \log \frac{E_i}{\sum_k E_k}, 0 \right) \quad (4.3)$$

(i and k run over the crystals in the supercluster) and where ω_0 is a constant that defines the energy threshold for a crystal to enter the sum.

4.1.2 Electron track reconstruction

The starting point of the track reconstruction is the seeding. Two approaches are used.

In the first one, named “ECAL driven” and designed for isolated electrons with transverse momentum above 5 GeV, the seed consists of the association of a supercluster with a transverse energy above 4 GeV with two compatible hits in the pixel and TEC layers. The supercluster energy and position are used to extrapolate the electron trajectory in the tracker and a hit compatible with the predicted position in a loose ϕ and z window is sought in the first or second layer of the pixel detector. Since the particle charge is unknown, two trajectories are built, corresponding to a +1 or -1 charge hypothesis. The position of the first hit is then used to constrain further the trajectory and a much narrower search window is used for the second hit search. The windows for the two hits depend on the tracker region and on the electron p_t . Low activity in the HCAL towers behind the supercluster is finally also required¹.

A “tracker driven” approach is also used and extends the electron track reconstruction for low momentum and non isolated electrons. It is developed in the context of the Particle Flow algorithm which aims to combine the information from the different subdetectors to reconstruct each particle produced in a collision individually and give a global description of the event. In this method, the starting point is a track and a search for a Particle Flow cluster is pursued. A multivariate analysis taking as input various track and cluster parameters is used to distinguish between the signal (electrons inside jets) from background (pions).

¹This takes the form of a condition on the H/E variable such as defined in section 4.3.3.

4. RECONSTRUCTION AND SELECTION OF HIGH ENERGY ELECTRONS WITH THE CMS DETECTOR

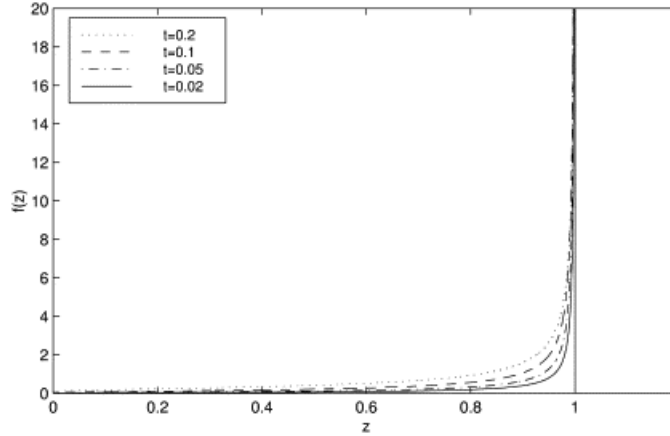


Figure 4.2: The Bethe-Heitler distribution modelling the remaining energy fraction carried by an electron interacting with matter. The distribution is shown for various fractions of radiation lengths, t [108].

4.1.2.1 Gaussian Sum Filter (GSF) tracks

As explained in section 2.3.1, the usual track reconstruction is made using a Kalman Filter algorithm. The assumption made that the energy loss in the tracker material follows a Gaussian distribution models well multiple scatterings but is not valid for electrons which follows a Bethe-Heitler distribution: a large amount of energy can sometimes be radiated in each thin layer of the tracker, inducing a non Gaussian large tail to the energy loss distribution, as illustrated in figure 4.2.

The hit collection is still built using a Kalman Filter, but with a very loose cut on their compatibility with the expected position and a more sophisticated algorithm, the Gaussian Sum Filter (GSF) algorithm [108, 109], is used to estimate the final track parameters. The idea is to approximate the Bethe-Heitler distribution distribution by a sum of Gaussians. The Gaussian Sum Filter is then a generalization of the Kalman Filter where the state vector in each tracker layer is described by a weighted sum of Kalman Filter tracks for which the energy loss follows a Gaussian law with a given width. The propagation of each mode is done separately from one site to another and the weights are then updated given the measurement in the new site. The practical consequence is that the allowed window to search for a hit in the next tracker layer is larger than for the usual Kalman Filter track.

4.1.3 GSF electron candidates

A GSF electron candidate is built by associating an ECAL supercluster with a GSF track with compatible η and ϕ positions. The variables used for the matching are defined in sections 4.3.1 and 4.3.2. Table 4.1 summarizes the set of conditions for a “ECAL driven” GSF electron candidate to be formed. The GSF electron candidate reconstruction efficiency is around 95% for electrons with a transverse momentum, p_t , above 35 GeV. Once a GSF electron candidate is constituted, the supercluster and track informations

Seeding: Two hits in the first layers of the tracker compatible with the trajectory extrapolated using the supercluster information. $H/E < 0.15$
Track-SC matching: $ \Delta\eta_{in} < 0.02$ $ \Delta\phi_{in} < 0.15$

Table 4.1: Set of conditions required to build a GSF electron candidate based on a ECAL driven seed.

are combined for the measurement of its momentum and charge.

For medium and low p_t values ($p_t < 50$ GeV), the momentum is obtained as a weighted sum of the supercluster energy and track momentum measurements. Electrons candidates are divided into different classes, depending on various criteria estimating the quality of the reconstruction and the amount of energy carried by bremsstrahlung photons. The weights, which depend on the electron class, η and p_t are derived from simulations using a regression technique. A clear improvement can be obtained for $p_t < 15$ GeV compared to the use of the supercluster information only. At transverse energy above 50 GeV however, the momentum measurement is completely driven by the supercluster. For the sake of simplicity and in order to avoid unexpected behaviour of the regression at very high p_t values, the electron momentum used in this thesis is therefore estimated using the supercluster energy only. The transverse momentum is therefore measured as

$$p_t = E_{SC} \sin \theta_{track} \quad (4.4)$$

where E_{SC} is the supercluster energy and θ_{track} the electron θ at the interaction point measured by the track. As the supercluster energy measurement is effectively used, the notation E_T rather than p_t will be used in the rest of this thesis.

Finally the charge of the candidate is estimated by combining three measurements: the GSF track curvature, the associated KF track curvature and the difference in ϕ between the position with respect to the beam spot of the supercluster and the first hit of the GSF track. The charge is given by the majority of the three measurements.

4.2 Electron triggers

As explained in section 2.2.7, the trigger system has two layers: the level 1 (L1) and the High Level Trigger (HLT).

At the level 1, where the tracker information is not available, electrons and photons are indistinguishable and based on calorimeter trigger towers, consisting, in the barrel, of a 5×5 matrix of ECAL crystals and the corresponding HCAL tower (a more complicated definition of the trigger tower is needed in the endcaps due to the different geometry). The L1 trigger starts from the tower with the highest ECAL energy deposit and $|\eta| < 2.5$, and builds a e/γ trigger candidate by combining it with the adjacent tower with the

4. RECONSTRUCTION AND SELECTION OF HIGH ENERGY ELECTRONS WITH THE CMS DETECTOR

highest energy. Two criteria on the shower shape are then applied: the ratio of the energy measured in the HCAL over the one in the ECAL is required to be less than 5% and, in the barrel, 90% of the ECAL energy must be contained in a 2×5 band of crystals in $\eta \times \phi$. A candidate is declared to be isolated if all the towers in a 3×3 matrix with an energy above a threshold also verify these criteria and if at least one of the 5 towers forming a corner of the matrix has a total energy below 3.5 GeV. The candidate is declared to be non isolated otherwise. The process is repeated until a maximum of 4 isolated and non isolated candidates are found. In the barrel, an additional condition is applied for a trigger tower to be considered. It was indeed noticed during the first LHC collisions the apparition of noise in the ECAL crystals due to the interaction of a low energy neutron directly in the avalanche photodiode of a crystal, leading to a high energy signal in this single crystal. In order to remove those “spikes”, and since electrons or photons leave at most around 80% of their energy in a single crystal, towers with an energy above a given threshold are requested to have at least two crystals with an energy above a given value. The full L1 trigger menu is a set of alternative conditions based on the various trigger objects. The main L1 conditions involving e/γ candidates require one candidate with typically $E_T > 20$ GeV or two candidates with E_T above around 10 GeV.

Events passing L1 are then filtered by the HLT. Here, the pixel tracker information is used to separate electrons from photons. The starting point of any electron HLT selection is a supercluster built as described in section 4.1.1. The trajectory of a charged particle with a position and energy given by the supercluster information is then extrapolated to the beam spot and a search for compatible hits in the pixel detector is performed. If at least two compatible hits are found, an electron candidate is constituted.

Many different triggers involving electrons are designed at the HLT level and various additional requirements on the electrons are made for each of them. They consist of conditions on:

- the longitudinal and transverse shower shape of the supercluster energy deposits.
- the ECAL and HCAL activity around the candidate.
- the existence of a KF track (or, rarely, a GSF track when the rate of passing events is low enough to allow the GSF algorithm be run) matching the supercluster position.
- the track activity around the candidate.

The conditions used and their severity depends on the number of electrons requested by the trigger and their E_T threshold, each trigger being designed to have a rate of passing events of 50 Hz or less.

Practically, all the HLT steps and criteria involving only calorimeters information are done first, the time consuming track reconstruction being only used at last for events passing the previous criteria.

The L1 and HLT trigger used to collect the data analyzed in this thesis are listed in table 4.2, which details a serie of conditions on several variables detailed in section 4.3. The L1 requires one e/γ candidate while the HLT requires two loose electrons candidates.

Name	Comments
L1_SingleEG22	1 e/γ candidate with $E_T > 22$ GeV
HLT_DoubleEle33_CaloIdL_GsfTrkIdVL	2 electrons objects satisfying: $E_T > 33$ GeV $ \Delta\eta_{in} < 0.02$; $ \Delta\phi_{in} < 0.15$ $H/E < 0.15$ (barrel) or 0.1 (endcaps) $\sigma_{in\eta} < 0.014$ (barrel) or 0.035 (endcaps)

Table 4.2: The L1 (first line) and HLT (second line) triggers used to collect the data analyzed in this thesis.

4.3 High E_T electron selection

The electron reconstruction presented in section 4.1 is the starting point of any electron selection in CMS and must be useable by any physics analysis involving one or several electrons in the final state. A high efficiency in any kinematical conditions is therefore needed. The consequence is that the probability for other particles such as jets to be reconstructed as electrons is sizeable.

In fact, at CMS, the GSF electron candidates reconstructed in the data are mainly fakes. These fakes are of several kinds: a charged pion can for instance mimic the signature of an electron if it interacts early and leaves most of its energy in the ECAL. Moreover electrons can emerge in a jet through the weakly decay of a hadron containing a c or b quark. Finally, in addition to jets, photons can also lead to GSF electron candidates. This happens if the photon converts into a dielectron pair in one of the first layers of the tracker detector. If one of the electron takes most of the photon momentum, a GSF electron candidate is likely to be reconstructed.

An analysis dependent selection, which takes into account the specific kinematics and background level needs therefore to be applied on top of the electron reconstruction. This thesis focuses on the search for very heavy resonances decaying into a dielectron pair, characterized by two energetic electrons in the final state. Most of the searches for new physics involving high energy electrons share the same features:

- The cross sections probed are tiny (sometimes below 1 fb).
- The background due to jets or other particles wrongly reconstructed as GSF electron candidates is moderate.
- The electron energy range under study is large (from $E_T = 100$ GeV to ≥ 1 TeV).
- The usual control region to check the performance of an electron selection in the data is the Z peak, where a large statistics is available for electrons with $E_T \approx 35$ GeV.
- For higher E_T values, the control regions (events with boosted on-shell Z or high mass tail of the Drell-Yan process) suffer from low statistics.

4. RECONSTRUCTION AND SELECTION OF HIGH ENERGY ELECTRONS WITH THE CMS DETECTOR

A high and stable selection efficiency for E_T above 100 GeV is therefore the main requirement. The selection must also reasonably work for lower values, down to 35 GeV. Finally, given the fact that the validation of the selection at high E_T in data is difficult, a simple and understandable selection is preferred.

In order to achieve this, a cut based selection has been developped [110]. It consists of a series of cuts on several variables that exploit the specificity of high E_T electrons. Only candidates with $E_T > 35$ GeV and well reconstructed in the tracker and ECAL sensitive regions are selected. Candidates in the ECAL transition region ($1.442 < |\eta_{SC}| < 1.56$, *cf.* section 2.2.5.1) and beyond the η coverage ($|\eta_{SC}| > 2.5$, *cf.* section 2.2.4) of the tracker are therefore discarded. A different selection is applied for candidates reconstructed in the ECAL barrel ($|\eta_{SC}| < 1.442$) and endcaps ($1.56 < |\eta_{SC}| < 2.5$).

The definition and the distribution of the various variables used are presented below for GSF electron candidates using the samples simulated at $\sqrt{s} = 8$ TeV presented in table 3.1, in the case of electrons (Drell-Yan process with $M_{ee} > 20$ GeV) and misidentified electrons (QCD or γ +jets events). In the Drell-Yan (DY) sample, the GSF electron candidates are required to match one of the generated electrons, there can therefore be two candidates per event. In the γ +jets sample, they are required to match the supercluster with the highest energy which, in more than 95% of the cases, coincides with the photon. In that purpose, one defines ΔR as:

$$\Delta R = \sqrt{\Delta\phi^2 + \Delta\eta^2} \quad (4.5)$$

where $\Delta\phi$ and $\Delta\eta$ are respectively the difference in ϕ and η between the GSF electron candidate and the generated electron (DY) or the supercluster (γ +jets) calculated at the interaction point. The matching condition is $\Delta R < 0.3$.

The cut value for each variable is given in table 4.3 and indicated in red in the figures. For comparison purposes, the cut values used for another high efficient electron selection, optimized for GSF electron candidates with $E_T \approx 35$ GeV and referred to as the “moderate E_T selection”, are also given.

The cut value on each variable was tuned such that it keeps around 99% of the GSF electron candidates with $E_T > 100$ GeV from the Drell-Yan process. Electron versus fake efficiency curves are presented in Appendix B for the main variables entering the selection. The Appendix B also shows the evolution of the efficiency as a function of E_T for each cut applied on GSF electron candidates (“single efficiency”) or on GSF electron candidates passing all the other cuts of the selection (“(N-1) efficiency”). The former allows one to quantify the intrinsic performance of a cut while the latter indicates how the various cuts are complementary. As will be seen below, most of these variables provide a good way to reject GSF electron candidates from jets: jets are expected to be composed of many particles and to interact with the calorimeters in a different way than electrons. The separation is however more difficult for converted photons since the candidate is here a well isolated electron. The shower shape and isolation variables are therefore expected to be of little help there.

4.3.1 $\Delta\eta_{in}$

This variable, the distribution of which is shown in figure 4.3, is defined as the difference in η , between the track position as measured in the inner layers, extrapolated to the interac-

Criteria	High E_T selection	Moderate E_T selection
SC-track matching	$ \Delta\eta_{in} < 0.005$ (0.007) $ \Delta\phi_{in} < 0.06$ (0.06)	$ \Delta\eta_{in} < 0.007$ (0.009) $ \Delta\phi_{in} < 0.15$ (0.1)
Longitudinal shower shape	$H/E < 0.05$ (0.05)	$H/E < 0.12$ (0.1)
Transverse shower shape	$E_{1\times5}/E_{5\times5} > 0.83$ OR $E_{2\times5}/E_{5\times5} > 0.94$ ($\sigma_{i\eta i\eta} < 0.03$)	$\sigma_{i\eta i\eta} < 0.01$ ($\sigma_{i\eta i\eta} < 0.03$)
Isolation	Tracker isolation < 5 GeV (5 GeV) AND ECAL+HCAL1 isolation $< 2+0.03E_T$ ($< \max(2.5, 1+0.03E_T)$)	PF Isolation $< 0.15E_T$ ($0.15E_T$)
Vertex compatibility	$ d_{xy} < 0.02$ (0.05)	$ d_{xy} < 0.02$ (0.02) $ d_z < 0.2$ (0.2)
Photon rejection	# missing hits ≤ 1 (1)	# missing hits ≤ 1 (1)

Table 4.3: The variables used in the high E_T and moderate E_T electron selections and the associated cut values for electrons emitted in the barrel (values in brackets indicate the corresponding cuts in the endcaps).

tion vertex and then extrapolated to the calorimeter and the position of the supercluster. The cut value in the barrel (0.005) is stronger than in the endcaps (0.007): the tracker material budget (*cf.* figure 2.8) is thicker in the latter case, reducing the precision on the position measurement. A much broader distribution is observed for fakes from QCD compared to electrons from DY. In the case of jets, all of its constituents can leave an energy deposit in the ECAL while the track only considers one of its constituents. The position of the center of the ECAL deposit can therefore be far from the track position.

4.3.2 $\Delta\phi_{in}$

$\Delta\phi_{in}$ (figure 4.4) is the analogous of $\Delta\eta_{in}$ for the ϕ coordinate. Its distribution is however much broader than $\Delta\eta_{in}$, because of the wider spread of the energy in ϕ with respect to η due to photons from bremsstrahlung. Hence, the cut value on $|\Delta\phi_{in}|$ (0.06 in both barrel and endcaps regions) is around ten times looser than for $\Delta\eta_{in}$. It is however tighter than the value used for the moderate E_T selection. This is because the distribution shrinks with increasing E_T as illustrated on figure 4.5 where the $|\Delta\phi_{in}|$ distribution is drawn as a function of E_T . Each constant E_T band is normalized to 1. The reason of this behaviour comes from the fact that bremsstrahlung photons, responsible for the main difference between the two measurements at moderate E_T are more colinear to the electron at higher E_T . One can be convinced that the bremsstrahlung photons are indeed the source of this effect by looking at the same distributions for Drell-Yan electrons and positrons separately (figure 4.6). One also sees in figure 4.4 the existence of two peaks in the $\Delta\phi_{in}$ distribution for the γ +jets events. The effect is easier to see for electrons reconstructed in the ECAL barrel but is also present for those reconstructed in the endcaps. It indicates a bias between the supercluster position and the track position. This can be explained by the following argument: in those events, a photon converts into two opposite sign electrons whose trajectory will bend in opposite ways in the ϕ direction, and one of the

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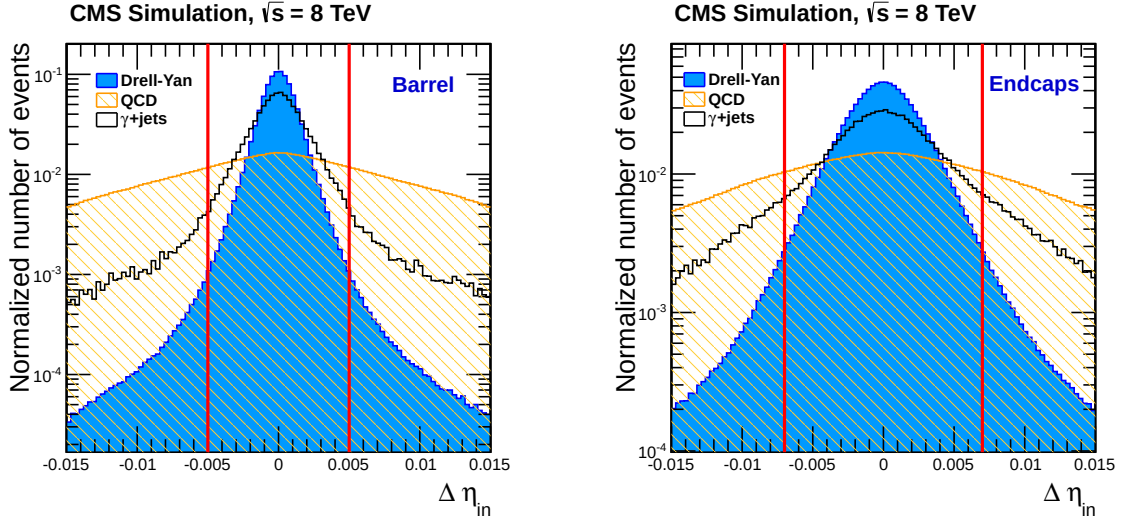


Figure 4.3: The distribution of $\Delta \eta_{in}$ for GSF electron candidates with $E_T > 35$ GeV emitted in the barrel (left) and in the endcaps (right) for Drell-Yan process (blue histogram), QCD events (dashed orange) and γ +jets events (black line). The red solid line indicates the cut value used in the high energy electron selection.

electron (usually the one with the highest energy) provides the track of the GSF electron candidate. Since the supercluster algorithm is designed to recover the energy of both electrons it will therefore have a position slightly shifted with respect to the position of the highest electron alone.

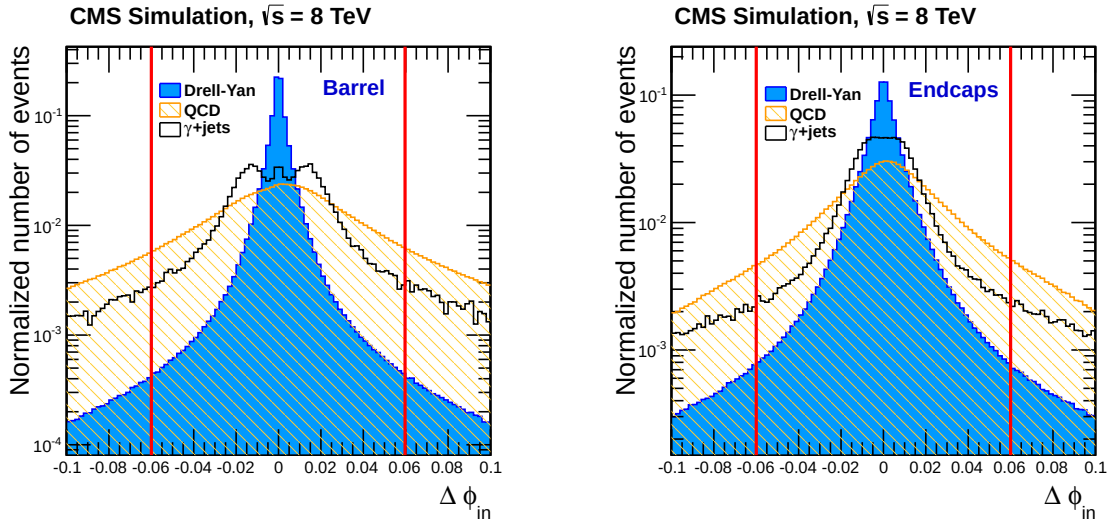


Figure 4.4: The distribution of $\Delta\phi_{in}$ for GSF electron candidates with $E_T > 35$ GeV emitted in the barrel (left) and in the endcaps (right) for Drell-Yan process (blue histogram), QCD events (dashed orange) and γ +jets events (black line). The red solid line indicates the cut value used in the high energy electron selection.

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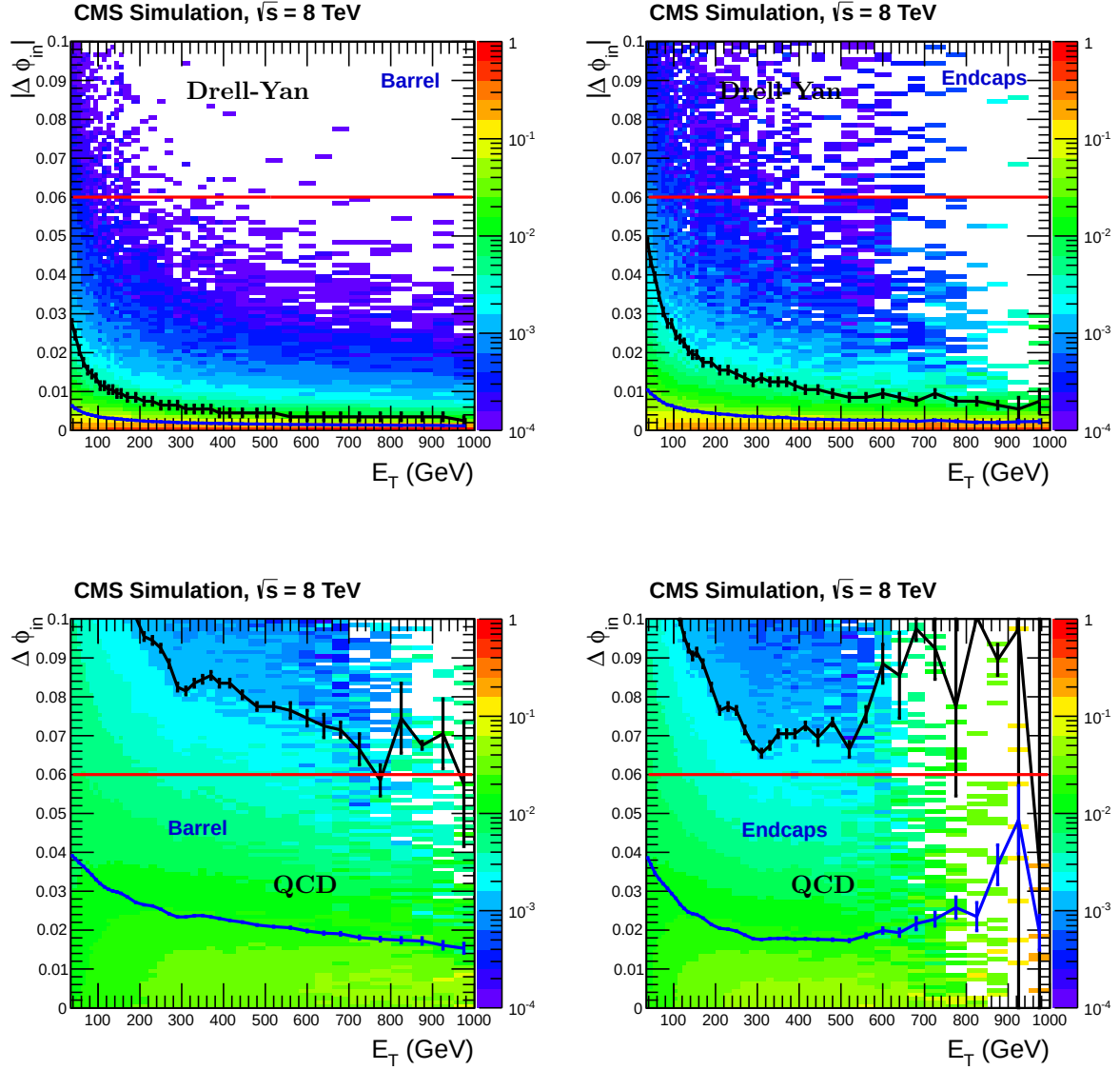


Figure 4.5: The distribution of $|\Delta\phi_{in}|$ for GSF electron candidates emitted in the barrel (left) and in the endcaps (right) as a function of E_T for Drell-Yan process (top) and QCD (bottom) events. Each E_T slice is normalized to 1 and for each of them, the blue and black curves show respectively the mean of the distribution and the value below which exactly 95% of the GSF electron candidates are found. The red solid line indicates the cut value used in the high energy electron selection.

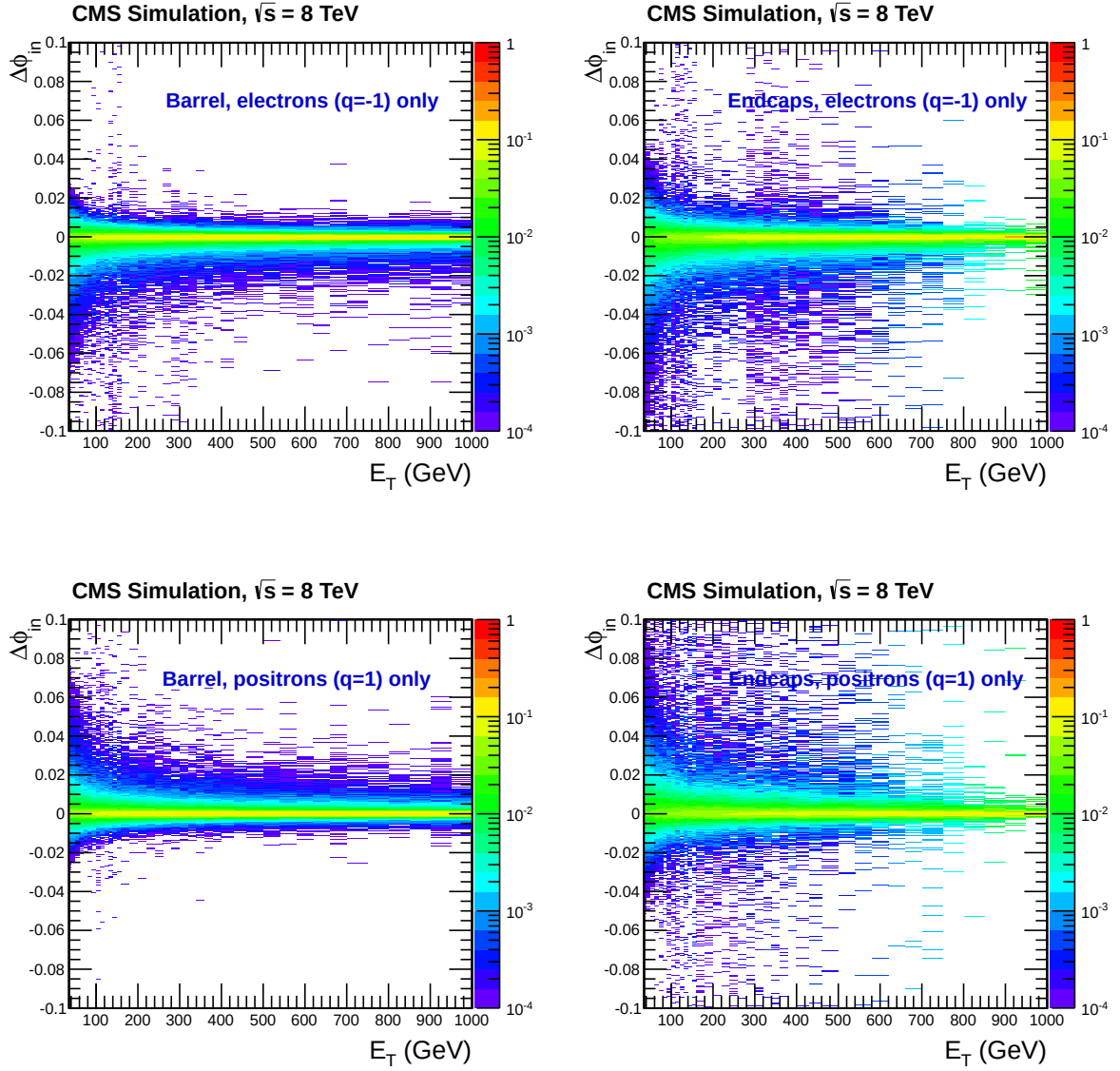


Figure 4.6: The distribution of $\Delta\phi_{in}$ for GSF electron candidates emitted in the barrel (left) and in the endcaps (right) as a function of E_T for electrons (top) and positrons (bottom) coming from the Drell-Yan process. Each E_T band is normalized to 1.

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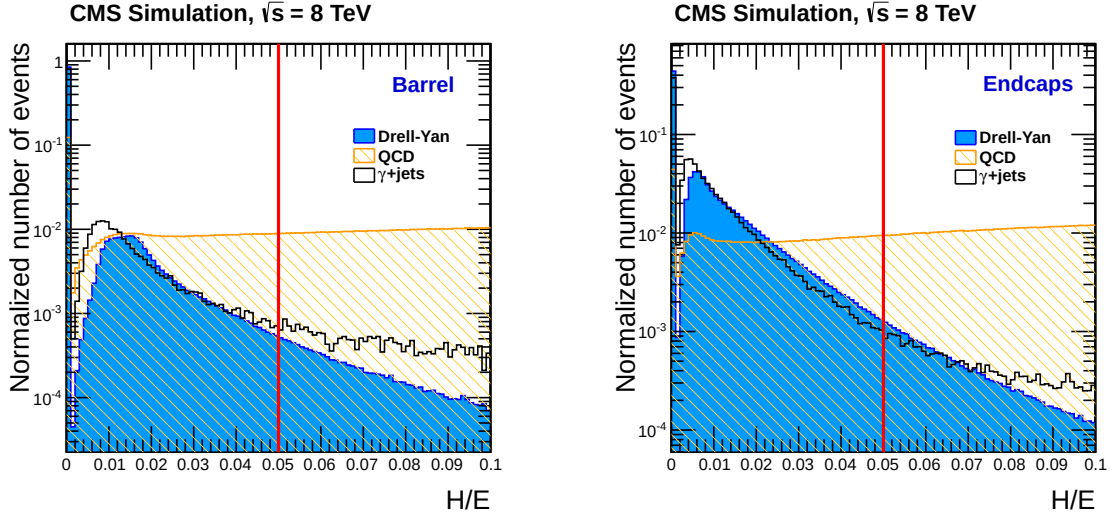


Figure 4.7: The distribution of H/E for GSF electron candidates with $E_T > 35$ GeV emitted in the barrel (left) and in the endcaps (right) for Drell-Yan process (blue histogram), QCD events (dashed orange) and γ +jets events (black line). The red solid line indicates the cut value used in the high energy electron selection.

4.3.3 H/E

H/E is defined as the ratio of the energy in the HCAL towers in a cone of radius $R = \sqrt{\Delta\phi^2 + \Delta\eta^2} = 0.15$ centered on the GSF electron candidate position to the electromagnetic energy of the GSF electron candidate supercluster. As for $\Delta\phi_{in}$, figures 4.7 and 4.8 show respectively the integrated distribution of H/E and its evolution as a function of the GSF electron candidate E_T . The cut value (0.05) is sensibly lower than the one usually used at moderate E_T (0.12 in the barrel and 0.10 in the endcaps). At moderate E_T this cut is indeed quite inefficient for a high number of pile-up interactions in the event, as will be discussed in section 4.4. However it is shown that the relative contribution of pile-up for electrons with E_T above 100 GeV is much smaller.

4.3.4 Transverse shower shape

Electrons or photons will leave energy in several ECAL crystals but their energy deposit distribution in the η coordinate is more narrow than other particles such as charged pions. This spread can be quantified through several variables:

- The ratio of the energy contained in the 1×5 domino in $\eta \times \phi$ in the barrel ($x \times y$ in the endcaps) centered on the seed crystal of the supercluster over the energy of the 5×5 matrix centered on the seed crystal, $\mathbf{E}_{1 \times 5} / \mathbf{E}_{5 \times 5}$.
- The ratio of the energy contained in the most energetic 2×5 domino in $\eta \times \phi$ in the

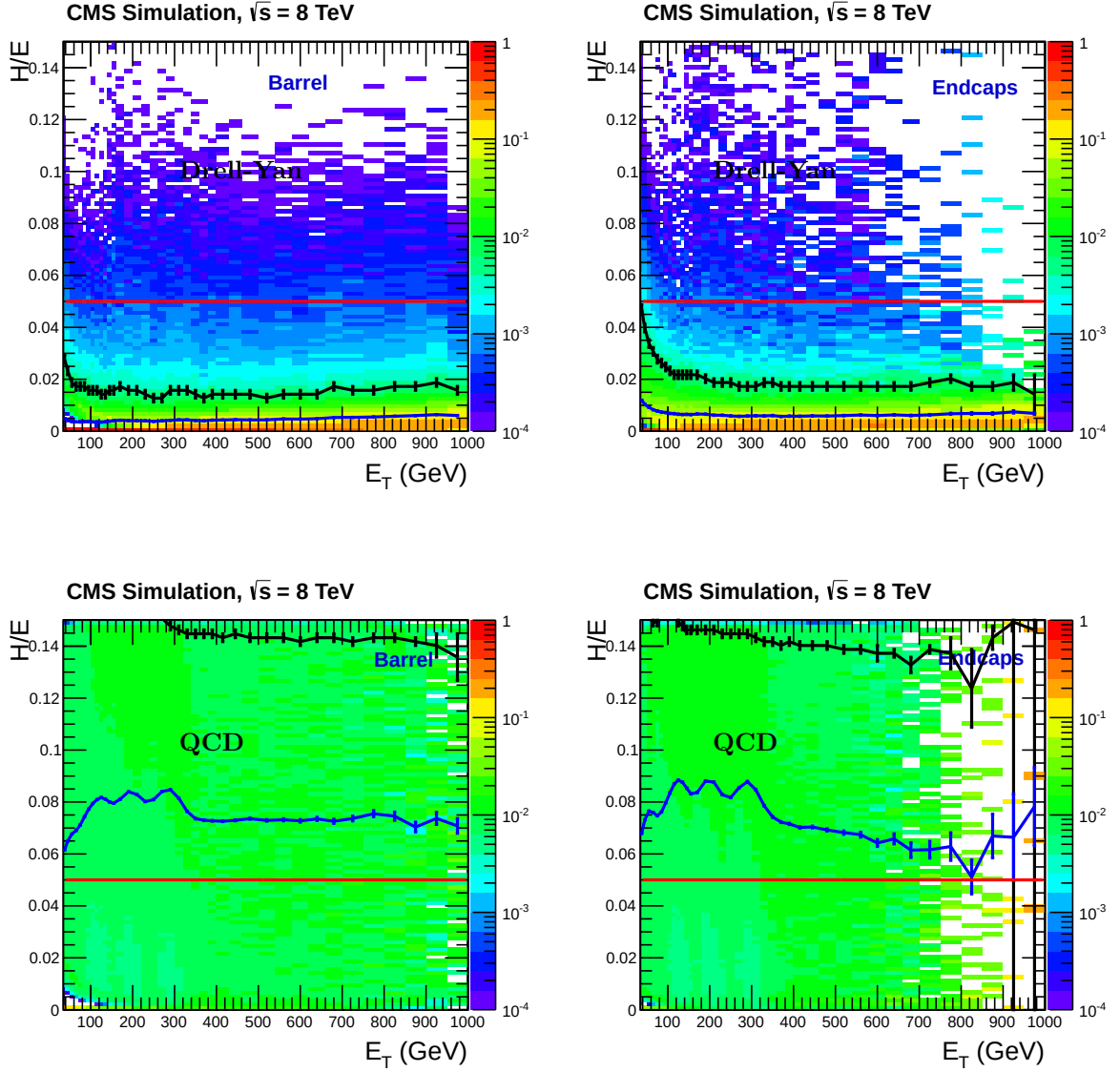


Figure 4.8: The distribution of H/E for GSF electron candidates emitted in the barrel (left) and in the endcaps (right) as a function of E_T for Drell-Yan process (top) and QCD (bottom) events. Each E_T slice is normalized to 1 and for each of them, the blue and black curves show respectively the mean of the distribution and the value below which exactly 95% of the GSF electron candidates are found. The red solid line indicates the cut value used in the high energy electron selection.

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barrel ($x \times y$ in the endcaps) centered on the seed crystal of the supercluster over the energy of the 5×5 matrix centered on the seed crystal, $\mathbf{E}_{2 \times 5}/\mathbf{E}_{5 \times 5}$.

- The η covariance, $\sigma_{i\eta i\eta}$, of the 5×5 matrix centered on the seed crystal:

$$\sigma_{i\eta i\eta}^2 = \frac{\sum_j^{5 \times 5} \omega_j \Delta \hat{\eta}_{crystal}^2 (i_{\eta,j} - i_{\eta,5 \times 5})^2}{\sum_j^{5 \times 5} \omega_j} \quad (4.6)$$

where:

- the weights ω_j are defined as in equation 4.3.
- $\Delta \hat{\eta}_{crystal}$ is the average crystal η width (0.01745 in the barrel and 0.0447 in the endcaps).
- $i_{\eta,j}$ is the signed distance, expressed in number of crystals, between the crystal j and the seed in the η direction¹.
- $i_{\eta,5 \times 5}$ is the signed distance, expressed in number of crystals, between the seed and the energy weighted position of the 5×5 matrix in the η direction.

The distributions of those three variables are shown in figures 4.9, 4.10 and 4.11.

The discrimination power of these variables is compared in figure 4.12 where the efficiency for electrons is plotted as a function of the efficiency for fake candidates from QCD events when applying a floating cut on $\sigma_{i\eta i\eta}$, on $E_{2 \times 5}/E_{5 \times 5}$ or on $E_{2 \times 5}/E_{5 \times 5}$ after a cut on $E_{1 \times 5}/E_{5 \times 5}$.

In the barrel, the best performance are obtained using a cut based on both $E_{1 \times 5}/E_{5 \times 5}$ and $E_{2 \times 5}/E_{5 \times 5}$. The two variables are indeed complementary: while $E_{1 \times 5}/E_{5 \times 5}$ is well designed for electrons hitting the center of a crystal, $E_{2 \times 5}/E_{5 \times 5}$ allows to recover electrons that hit the crystal close to its edge. Combining the two variables instead of using just one of them allows to set strong cut values on both and thus well reject background while keeping a high efficiency on simulated electrons. This feature is illustrated on the scatter plot in figure 4.13 for both barrel and endcaps electrons.

One sees in figure 4.13 that the distribution is much broader for signal electrons in the endcaps. This is because the $x - y$ geometry does not follow the $\eta - \phi$ direction. As a consequence, part of the energy is sometimes missing in both $E_{1 \times 5}/E_{5 \times 5}$ and $E_{2 \times 5}/E_{5 \times 5}$ and their discrimination power is reduced and becomes worse than $\sigma_{i\eta i\eta}$.

4.3.5 Missing hits in the innermost layers of the trackers

This variable is defined as the number of missing hits in the innermost layers of the tracker (including the pixel) before the GSF track first hit. It is mainly designed to reject photons that convert into a pair in the tracker as can be seen from figure 4.14. It can happen that only one of the electron tracks is reconstructed. The signature left in such a case is then

¹ This definition and the next one is straightforward in the barrel given that the crystals are arranged in constant η or ϕ bands. In the endcaps, the crystals are arranged in a $x-y$ grid centered on the z -axis. Denoting x_j and y_j the coordinates, expressed in number of crystal width, of the crystal j in the grid, one defines R_j as $\sqrt{x_j^2 + y_j^2}$ and $i_{\eta,j}$ is then defined as $R_j - R_{seed}$.

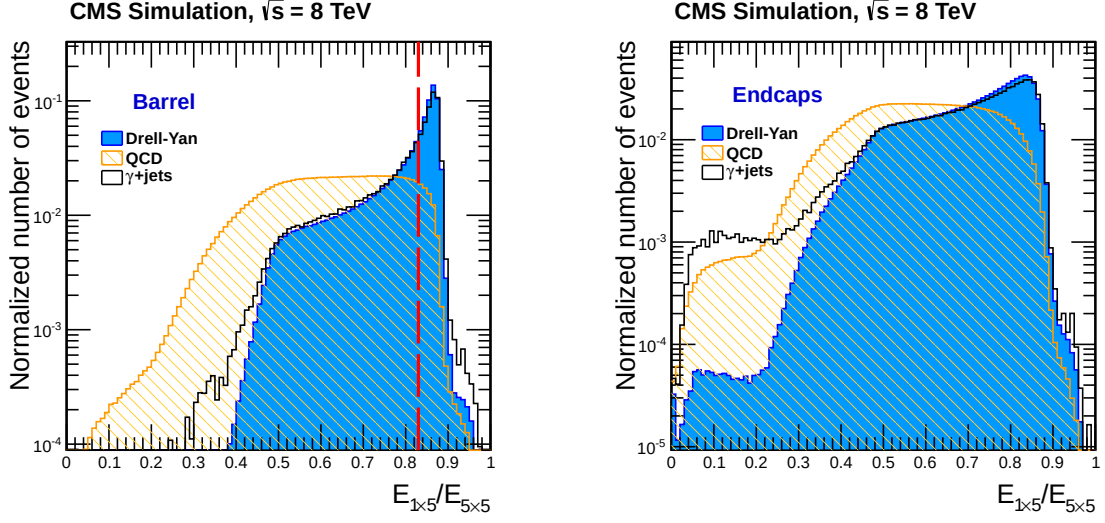


Figure 4.9: The distribution of $E_{1\times 5}/E_{5\times 5}$ for GSF electron candidates with $E_T > 35$ GeV emitted in the barrel (left) and in the endcaps (right) for Drell-Yan process (blue histogram), QCD events (dashed orange) and γ +jets events (black line). This variable is used in combination with $E_{2\times 5}/E_{5\times 5}$ in the high energy electron selection for barrel electrons which are requested to pass the cut on one of the two variables. The red dashed line indicates the cut value for $E_{1\times 5}/E_{5\times 5}$.

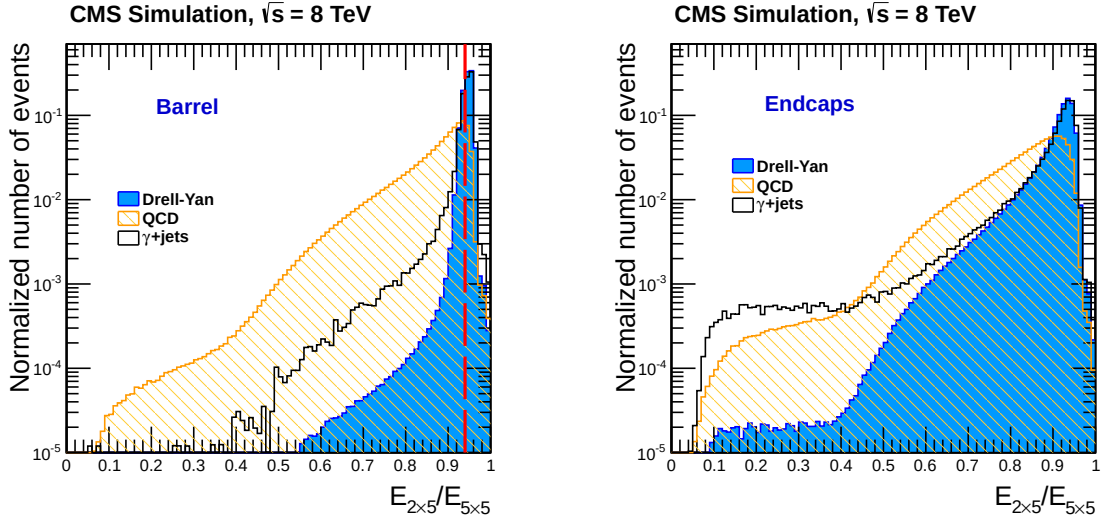


Figure 4.10: The distribution of $E_{2\times 5}/E_{5\times 5}$ for GSF electron candidates with $E_T > 35$ GeV emitted in the barrel (left) and in the endcaps (right) for Drell-Yan process (blue histogram), QCD events (dashed orange) and γ +jets events (black line). This variable is used in combination with $E_{1\times 5}/E_{5\times 5}$ in the high energy electron selection for barrel electrons which are requested to pass the cut on one of the two variables. The dashed line indicates the cut value for $E_{2\times 5}/E_{5\times 5}$.

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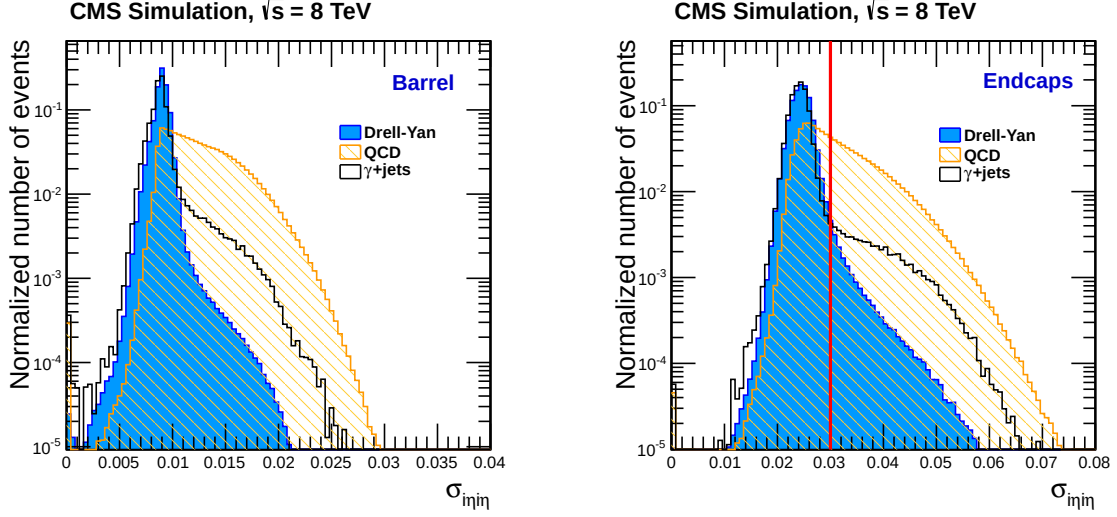


Figure 4.11: The distribution of σ_{inj} for GSF electron candidates with $E_T > 35$ GeV emitted in the barrel (left) and in the endcaps (right) for Drell-Yan process (blue histogram), QCD events (dashed orange) and γ +jets events (black line). The red solid line indicates the cut value used in the high energy electron selection for endcaps electrons.

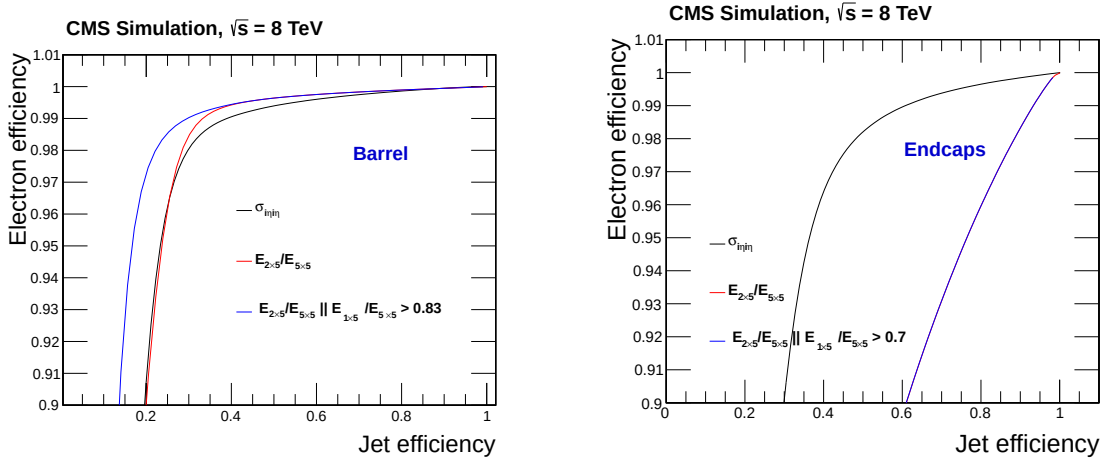


Figure 4.12: The efficiency of a cut on various lateral shower shape variables σ_{inj} , $E_{2\times 5}/E_{5\times 5}$ and $E_{2\times 5}/E_{5\times 5}$ with the alternative condition $E_{1\times 5}/E_{5\times 5} > 0.83$, for GSF electron candidates emitted in the barrel (left) or endcaps (right), for simulated electrons from the Drell-Yan process as a function of the same efficiency for GSF electron candidates from fake electrons in QCD events.

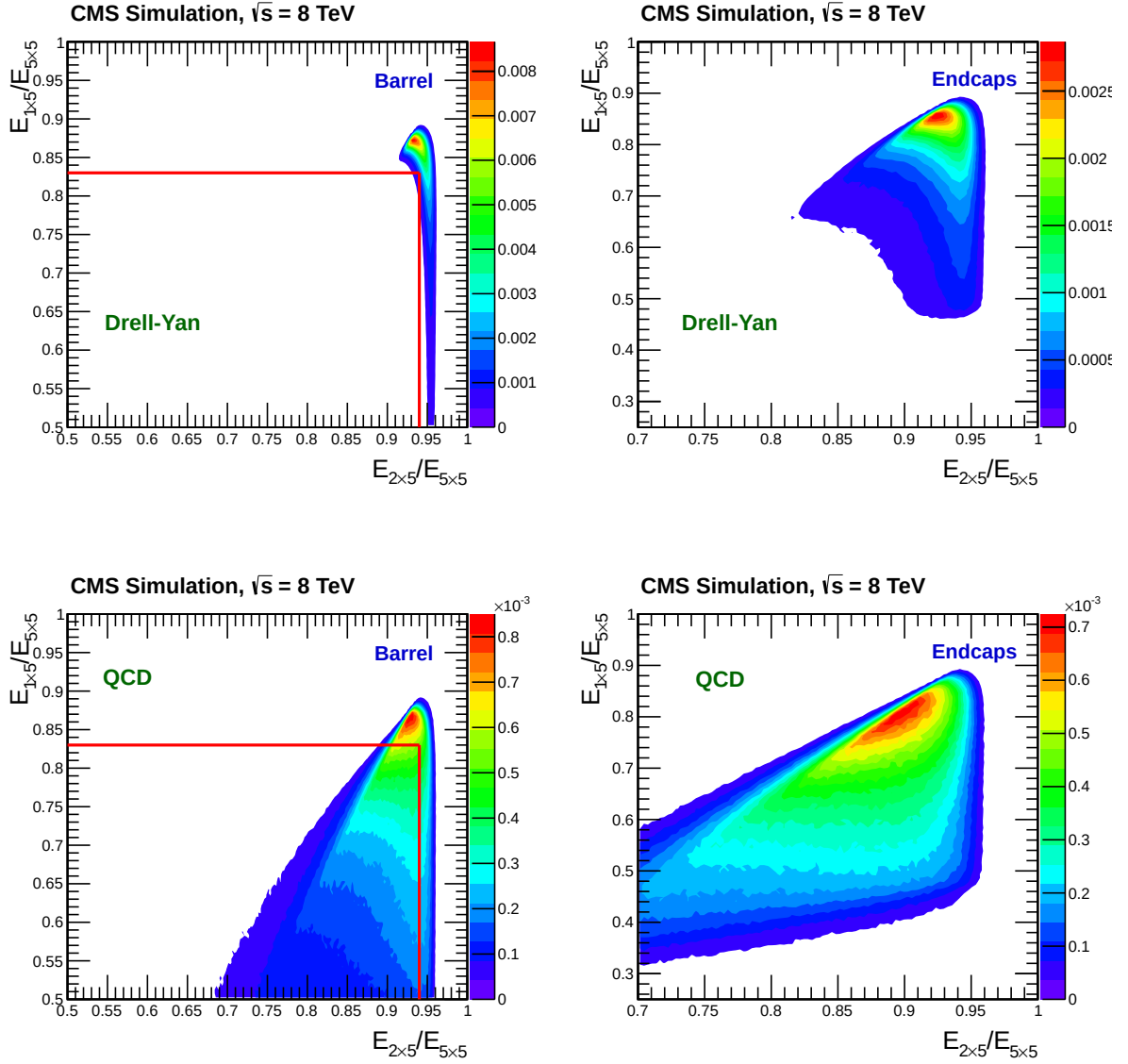


Figure 4.13: The distribution of E_{1x5}/E_{5x5} vs E_{2x5}/E_{5x5} for GSF electron candidates with $E_T > 35$ GeV emitted in the barrel (left) and in the endcaps (right) for Drell-Yan process (top) and QCD (bottom) events. The red solid lines indicate the cut values used in the high energy electron selection for barrel electrons.

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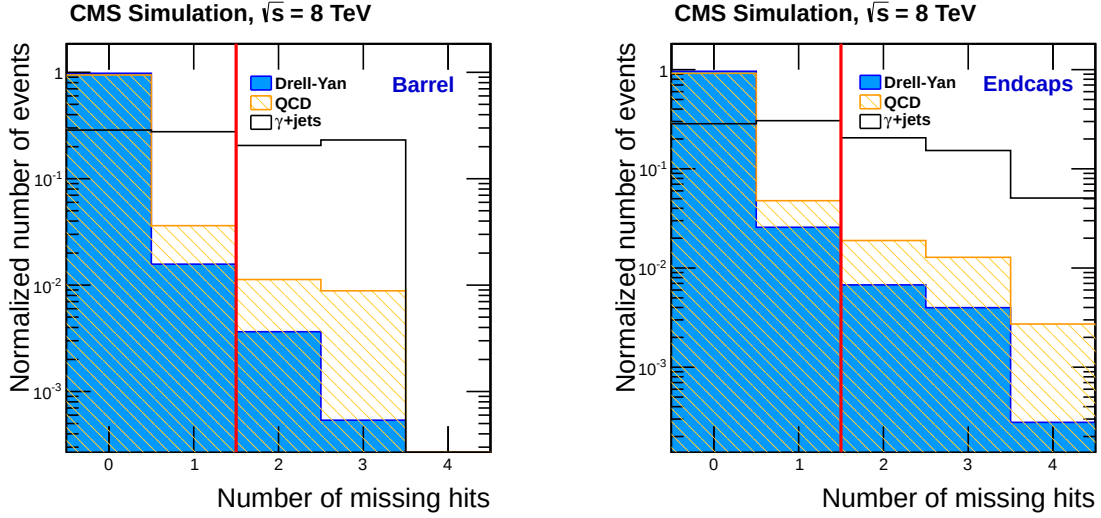


Figure 4.14: The distribution of the number of missing hits in the tracker inner layers for GSF electron candidates with $E_T > 35$ GeV emitted in the barrel (left) and in the endcaps (right) for Drell-Yan process (blue histogram), QCD events (dashed orange) and γ +jets events (black line). The red solid line indicates the cut value used in the high energy electron selection.

very similar to the one from real electrons. In particular, the shower shape variables are of no help to disentangle them. One of the main differences however is the absence of hits in the first layers of the tracker, before the conversion happens.

4.3.6 Transverse impact parameter (d_{xy})

The transverse impact parameter (figure 4.15) is the closest distance, in the transverse plane, between the primary vertex and the track of the GSF electron candidate. The distribution is wider in the endcaps due to the poorer resolution on the track momentum in that region. Similarly to the missing hit cut, the d_{xy} cut is mainly useful to reject converted photons.

4.3.7 Calorimeter isolation ECAL+HCAL1

The calorimeter isolation variable used in the high E_T selection is the sum of:

- **The ECAL isolation** defined as the transverse energy of all the crystals with $E > 80$ MeV in the barrel ($E_T > 100$ MeV in the endcaps) in a cone of $\Delta R = 0.3$ centered on the GSF electron candidate position in the calorimeter excluding those in an inner cone of radius 3 crystals and η strip of total width of 3 crystals (*cf.* figure 4.16 left). The inner cone veto removes the electron energy from the sum whereas the η strip removes bremsstrahlung photons.

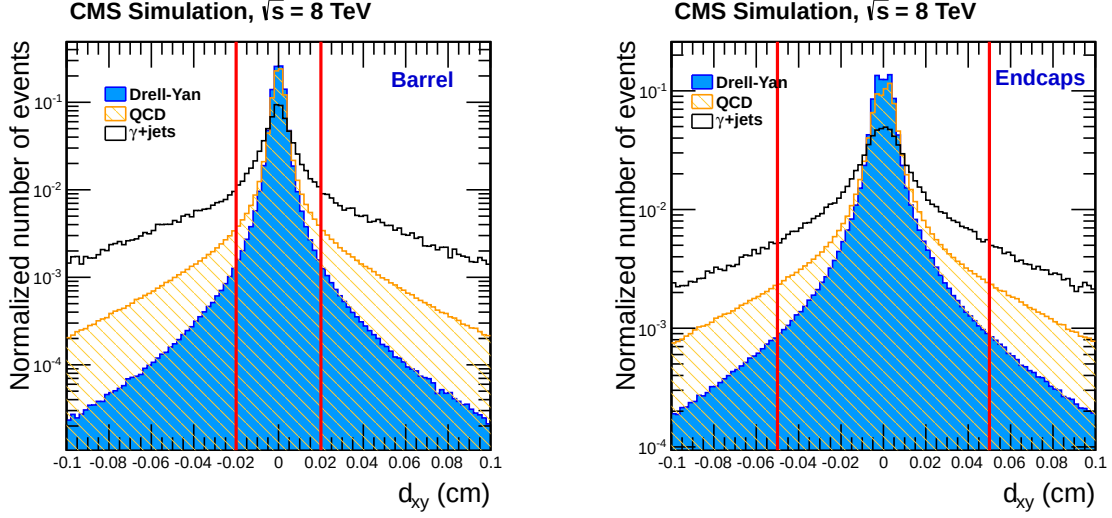


Figure 4.15: The distribution of d_{xy} for GSF electron candidates with $E_T > 35$ GeV emitted in the barrel (left) and in the endcaps (right) for Drell-Yan process (blue histogram), QCD events (dashed orange) and γ +jets events (black line). The red solid line indicates the cut value used in the high energy electron selection.

- **The HCAL1 isolation** defined as the transverse energy collected by all the towers of the first layer of the HCAL in a cone of $\Delta R = 0.3$ centered on the GSF electron candidate position in the calorimeter, excluding the towers in a cone of $\Delta R = 0.15$ (*cf.* figure 4.16 right). The first depth is defined as all depths for towers 1 to 17, depth 1 for towers 18 to 29 and depth 1 and 2 for towers 27 to 29 (*cf.* section 2.2.5.2).

The distribution of the ECAL+HCAL1 isolation is shown in figure 4.17 and its evolution with E_T in figure 4.18. The variable is very dependent of the electron energy and tends to increase for high E_T electrons due to the extension of the shower. A cut value of the form $Iso < A + B \times E_T$ is therefore applied. The parameter A has different values for the barrel and endcaps but the parameter B is identical in both regions and is 3%.

This cut differs from the moderate E_T selection in two aspects. First, instead of an isolation using separately each subdetectors information, moderate and low E_T selections rather use a Particle Flow (PF) based isolation which brings together the information of the tracker, the ECAL and the HCAL which allows to measure the contribution to the isolation from “charged hadrons”, “neutral hadrons” and “photons” separately. One of the main advantage of proceeding so is that the energy deposit in the calorimeters associated to a charged hadron produced in another interaction, characterized by a different primary vertex, can be removed from the isolation sum. For very high energy (>1 TeV) electrons, however, the PF sometimes fails to recognize an electron from a GSF electron candidate and assigns all its energy deposit to the photon isolation. Secondly, the combined PF isolation is generally requested to be below 15% of the electron E_T . This is perfectly acceptable at moderate E_T ($E_T < 50$ -100 GeV) but this would mean a cut value of

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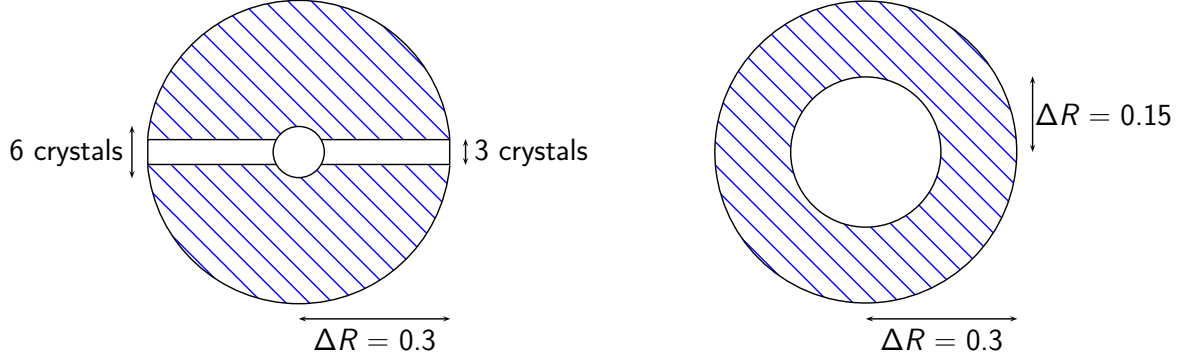


Figure 4.16: Definition of the regions used to compute the ECAL (left) and HCAL1 (right) isolation (shaded area). The disks are centered on the electron position.

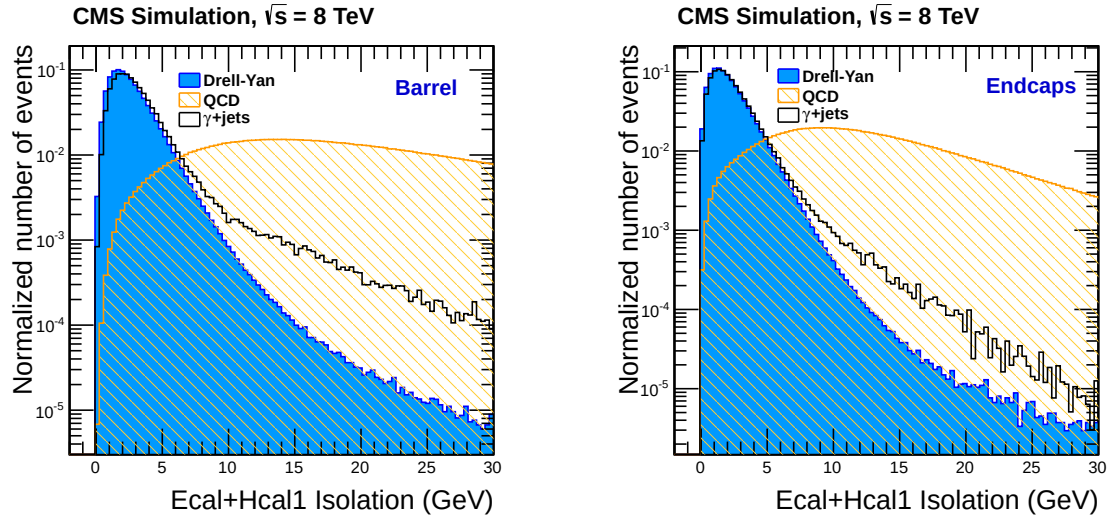


Figure 4.17: The distribution of the ECAL+HCAL1 isolation for GSF electron candidates with $E_T > 35$ GeV emitted in the barrel (left) and in the endcaps (right) for Drell-Yan process (blue histogram), QCD events (dashed orange) and γ +jets events (black line).

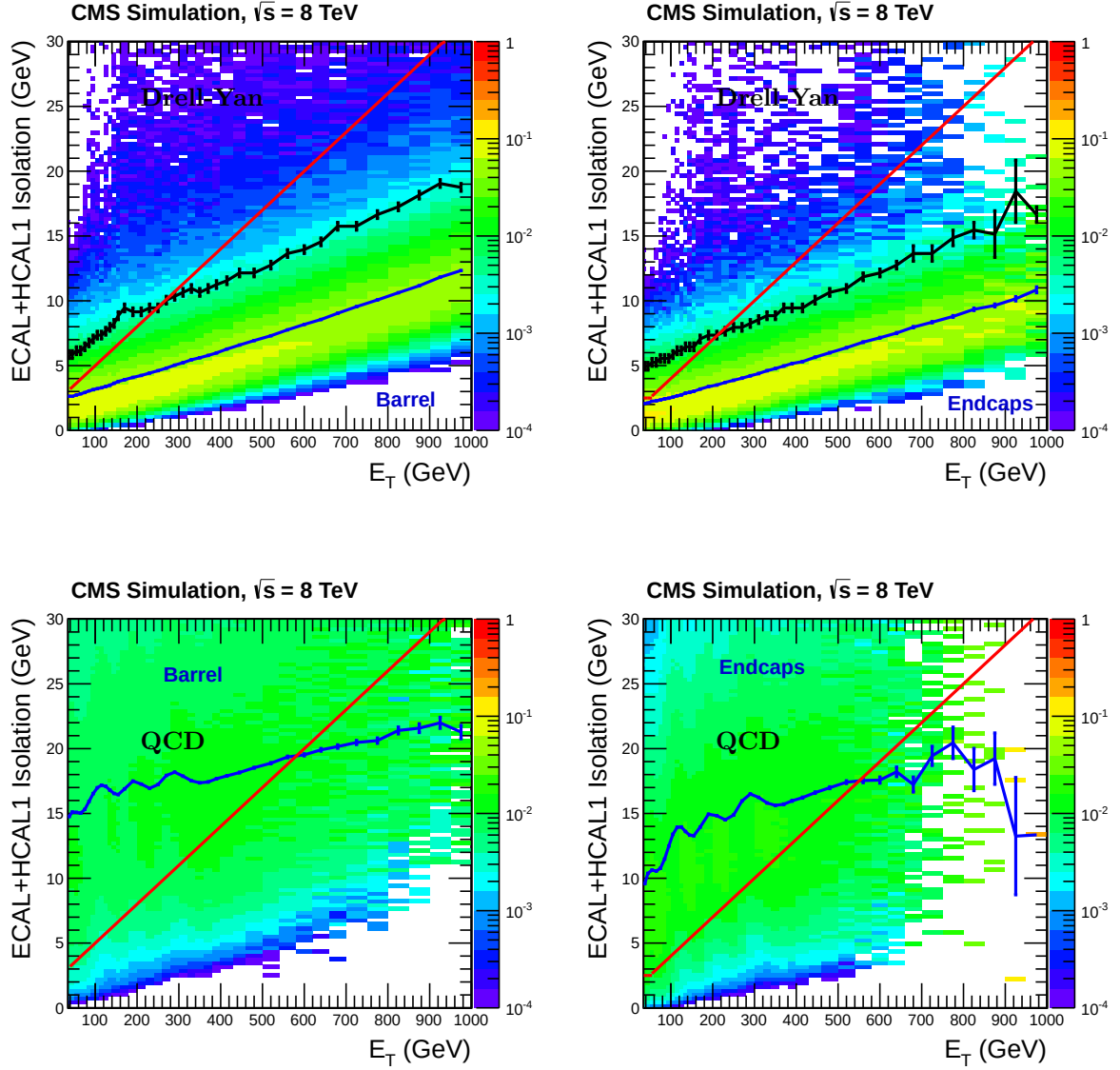


Figure 4.18: The distribution of the ECAL+HCAL1 Isolation for GSF electron candidates emitted in the barrel (left) and in the endcaps (right) as a function of E_T for Drell-Yan process (top) and QCD (bottom) events. Each E_T slice is normalized to 1 and for each of them, the blue and black curves show respectively the mean of the distribution and the value below which exactly 95% of the GSF electron candidates are found. The red solid line indicates the cut value used in the high energy electron selection.

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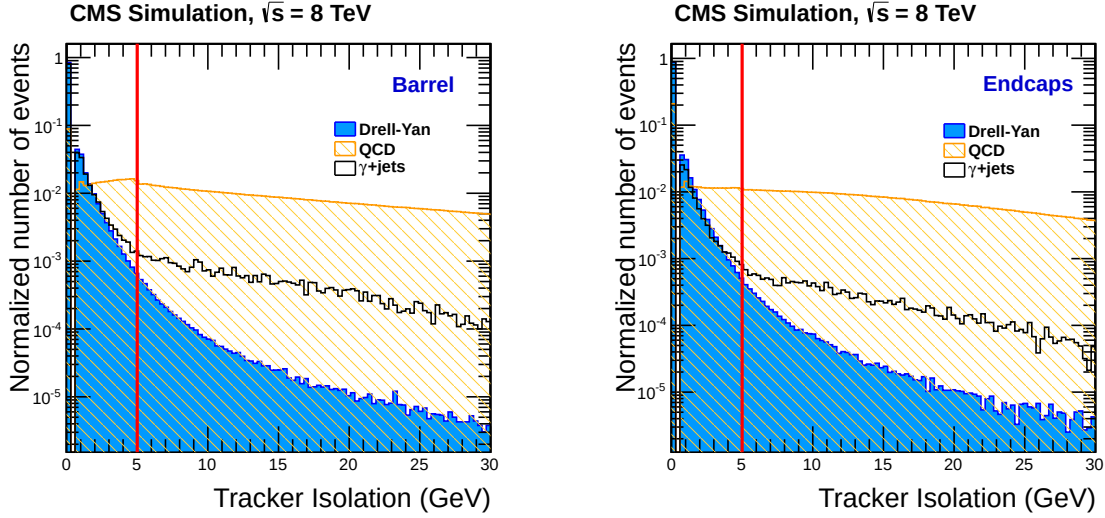


Figure 4.19: The distribution of the track isolation for GSF electron candidates with $E_T > 35$ GeV emitted in the barrel (left) and in the endcaps (right) for Drell-Yan process (blue histogram), QCD events (dashed orange) and γ +jets events (black line). The red solid line indicates the cut value used in the high energy electron selection.

150 GeV for an electron with $E_T = 1$ TeV which is far above the typical values for real electrons.

4.3.8 Tracker isolation

This variable (*cf.* figure 4.19) is defined as the sum of the p_t of all the reconstructed KF tracks between an inner and outer cone defined respectively as $\Delta R=0.04$ and 0.3 and centered on the GSF electron candidate direction. Only tracks having $p_t > 700$ MeV and $|\Delta z|$ with the electron track < 0.2 are considered, where z is the minimum distance of the track to (0,0,0). This variable presents the advantage of being very little sensitive to the pile-up (tracks from other vertices are rejected) and does not show a dependency with the electron E_T , as can be concluded from 4.20. The inner cone veto is presents to remove duplicate tracks associated to the electron as well as tracks from converted bremsstrahlung photons.

4.3.9 E/p

The ratio of the supercluster energy over the momentum as measured by the track, E/p , is another variable that can help disentangling electrons from fakes. A distribution peaking at 1 is indeed expected in the former case whereas the two measurements are only weakly correlated for fakes. This behaviour is indeed observed in figure 4.21. However a large tail, especially in the endcaps is present for electrons. This comes from the fact that the track

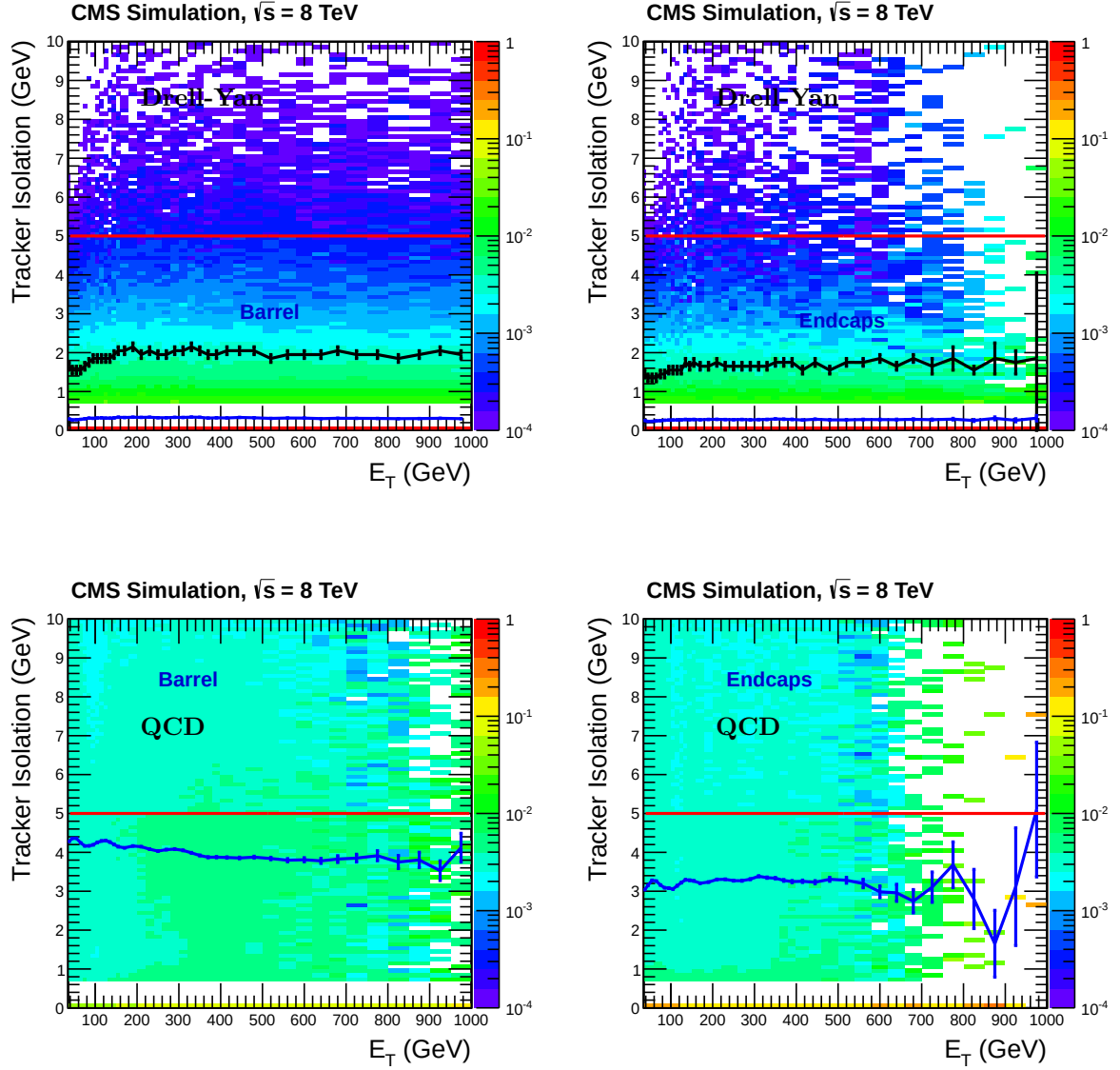


Figure 4.20: The distribution of the track isolation for GSF electron candidates emitted in the barrel (left) and in the endcaps (right) as a function of E_T for Drell-Yan process (top) and QCD (bottom) events. Each E_T slice is normalized to 1 and for each of them, the blue and black curves show respectively the mean of the distribution and the value below which exactly 95% of the GSF electron candidates are found. The red solid line indicates the cut value used in the high energy electron selection.

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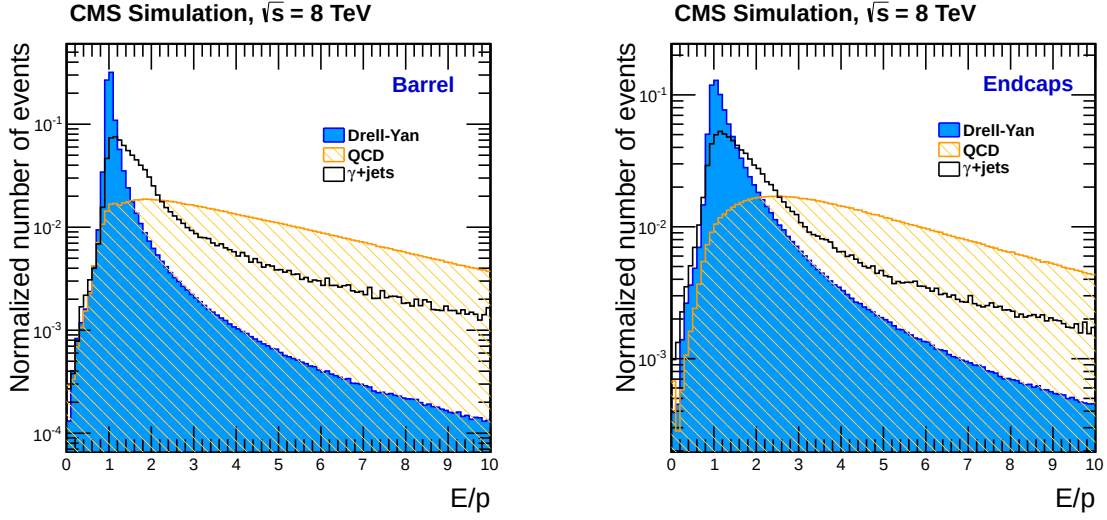


Figure 4.21: The distribution of E/p for GSF electron candidates with $E_T > 35$ GeV emitted in the barrel (left) and in the endcaps (right) for Drell-Yan process (blue histogram), QCD events (dashed orange) and γ +jets events (black line).

momentum measurement is subject to a large uncertainty due to the small curvature of high E_T particles in the detector. This is more clear when looking at the E/p evolution with E_T in figure 4.22. The variable can help to distinguish between electrons and fakes but for a high efficiency selection, a cut on this variable is impossible.

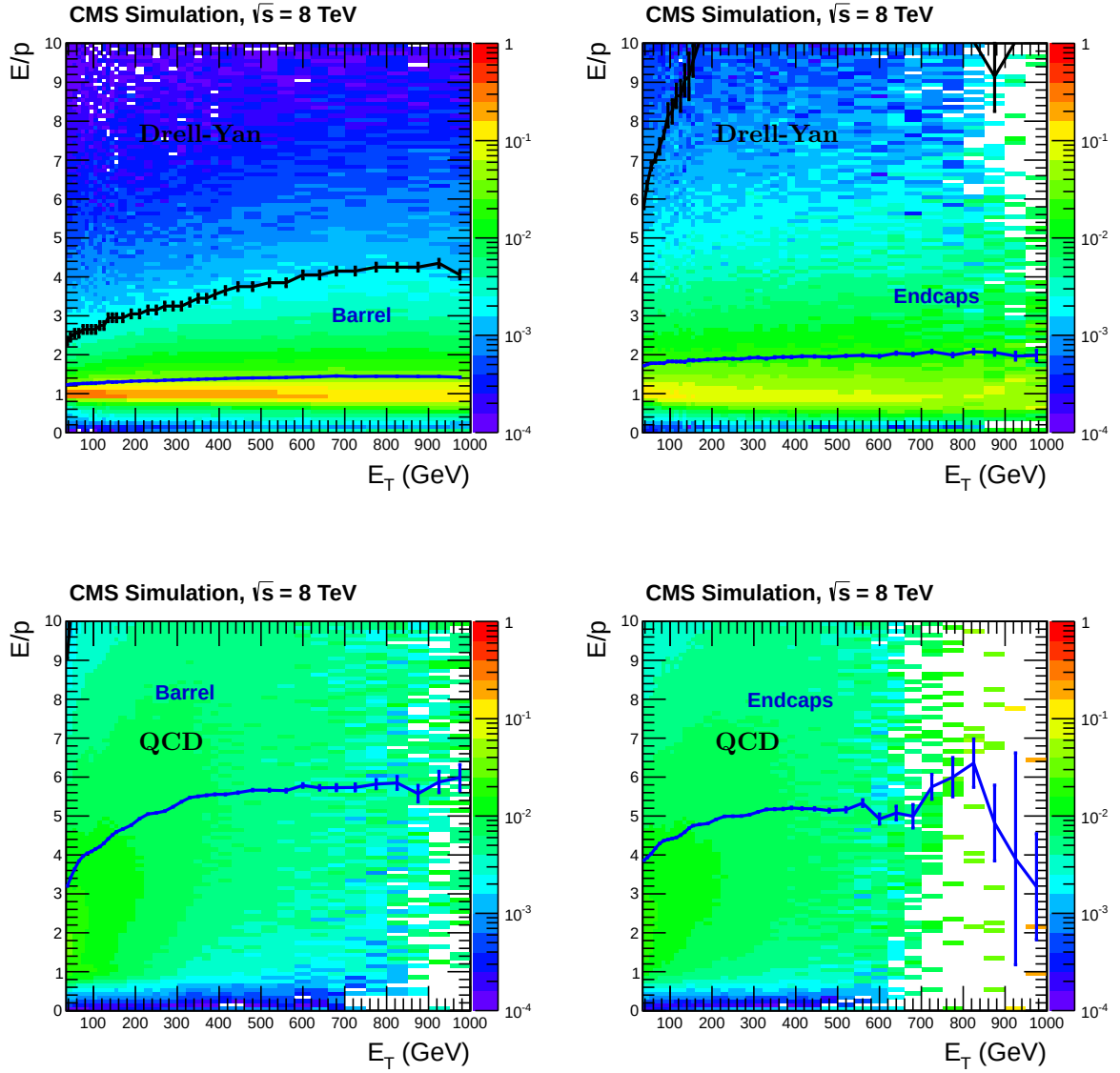


Figure 4.22: The distribution of E/p for GSF electron candidates emitted in the Barrel (left) and in the Endcaps (right) as a function of E_T for Drell-Yan process (top) and QCD (bottom) events. Each E_T band is normalized to 1.

4.4 Pile-up corrections

As discussed in section 2.1.3, at the LHC, several collisions occur in the same bunch crossing. In addition to the particles coming from the hard scattering collision under study, particles from the piled up collisions are also measured in the detector and can potentially alter the event description and affect the objects reconstruction. This section studies the effect of pile-up events on the electron selection. It is shown that several of the variables mentioned in the previous section are affected. In particular, the amount of energy in the calorimeter isolation cone is strongly sensitive to pile-up leading to a dramatic drop of the efficiency for large number of pile-up interactions. A technique using the average transverse energy density due to pile-up in the event to subtract this contribution is therefore studied. The results presented in this section were initially obtained with simulations at $\sqrt{s} = 7$ TeV [110] and were reproduced with the latest simulations available at $\sqrt{s} = 8$ TeV for this thesis only. It is therefore the corrections derived with the former simulations that were used practically.

4.4.1 Efficiency before pile-up correction

Figure 4.23 presents the total high E_T electron selection efficiency (starting from a GSF electron candidate) as a function of the number of reconstructed primary vertices N_{vtx} for Drell-Yan simulated events, separately for barrel and endcaps electrons with $E_T > 35$ GeV (top) or 100 GeV (bottom). A clear dependence is observed as function of N_{vtx} , especially at moderate E_T . The origin of this dependency can be found by looking at figures 4.24, 4.25 and 4.26 where the efficiency of each cut entering the selection is presented separately. A sizeable reduction of the efficiency at high number of reconstructed vertices is observed for three cuts: the ECAL+HCAL1 isolation, the transverse and longitudinal shower shape variables, the most affected being the ECAL+HCAL1 isolation one.

4.4.2 The ρ variable and the “effective area” method

In order to restore a high efficiency of the selection in high pile-up conditions, a technique to subtract the pile-up noise has been developped. It was initially studied to correct the calorimeter isolation criteria, as these are the most strongly affected. The principle is to subtract the estimated energy contamination due to pile-up from the isolation sum. This is achieved using the ρ variable defined as the transverse energy density due to the pile-up interactions. The variable can be calculated for each event using the Fastjet framework [111, 112, 113] which proposes various techniques to find and reconstruct jets. Once all the particles and jets associated to the hard interaction are identified, one defines “quiet” regions in the event. The energy density in those regions is defined as the pile-up energy density. At $\sqrt{s} = 8$ TeV, ρ was measured to be of the order 1 GeV/ $(\Delta\eta\Delta\phi)$ per vertex. The corrected isolation, Iso_{corr} , is obtained as:

$$Iso_{corr} = Iso_{raw} - \rho \times E_A \quad (4.7)$$

where Iso_{raw} is the isolation value before correction and E_A represents an effective area. Instead of using a geometric value for E_A , the effective area is computed from the slopes

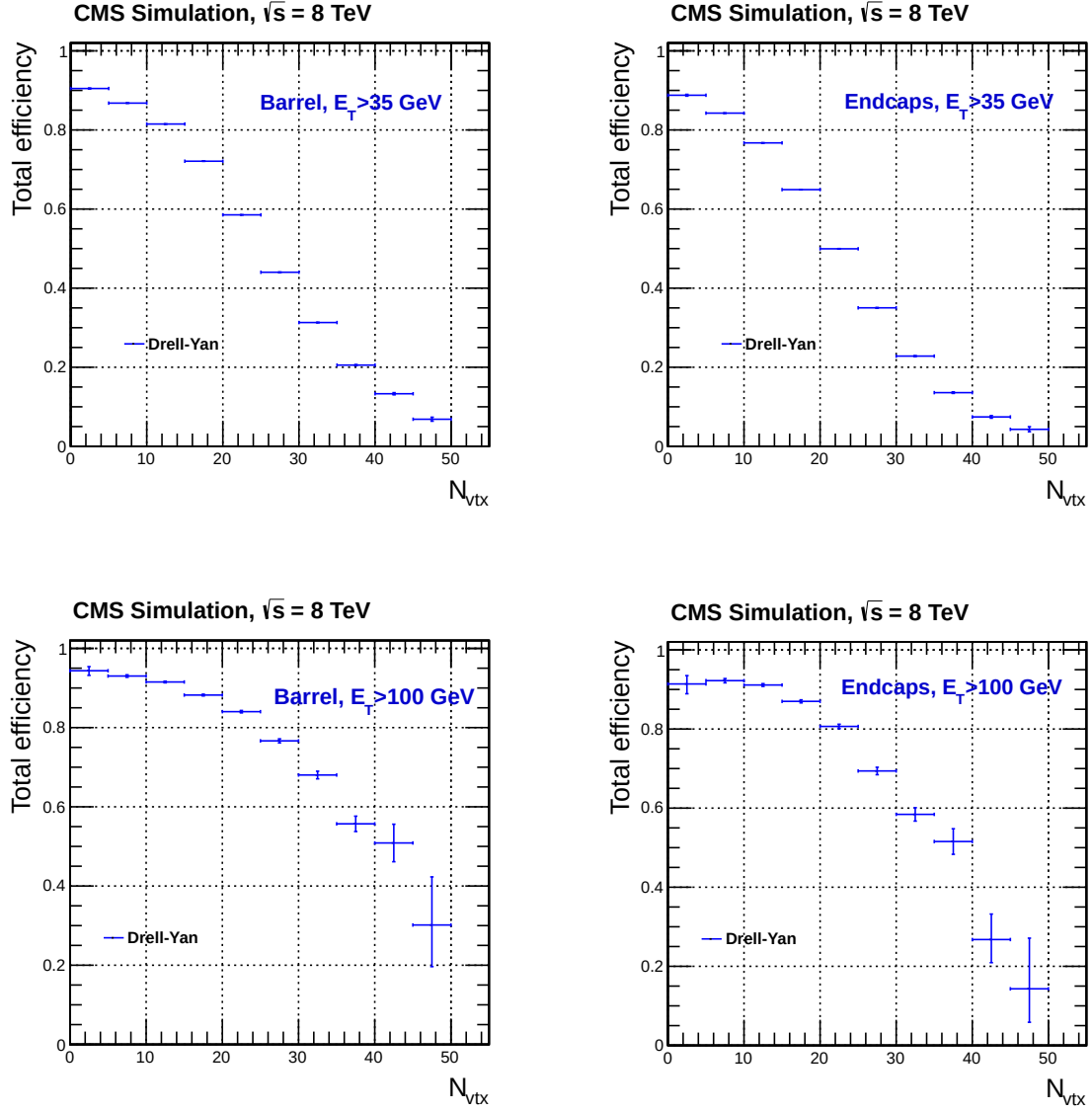
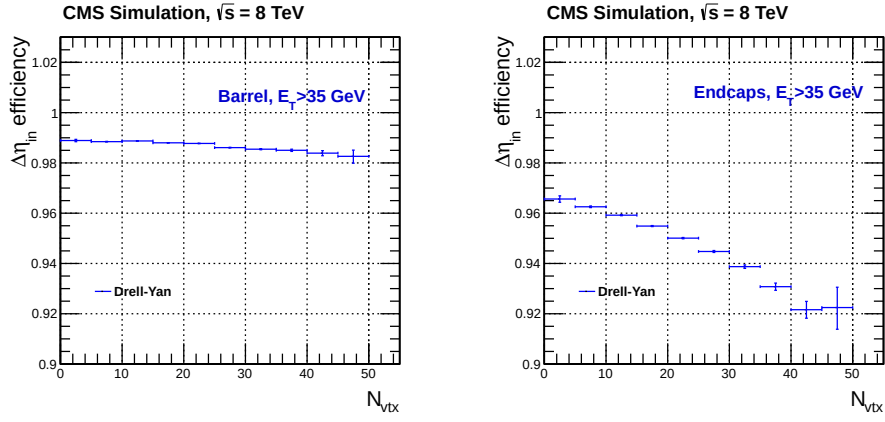
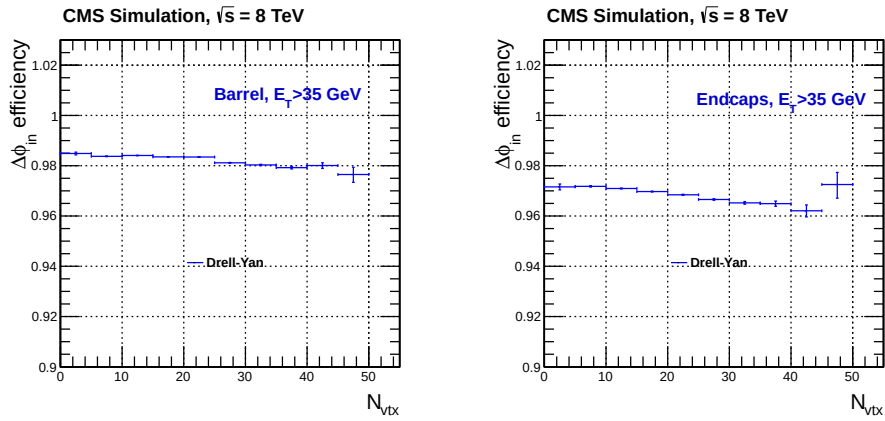


Figure 4.23: Total high E_T electron selection efficiency starting from a GSF electron candidate as a function of the number of vertices, N_{vtx} , for Drell-Yan simulated events, separately in the ECAL barrel (left) and endcaps (right), for electrons with $E_T > 35$ GeV (top) and $E_T > 100$ (bottom) GeV.

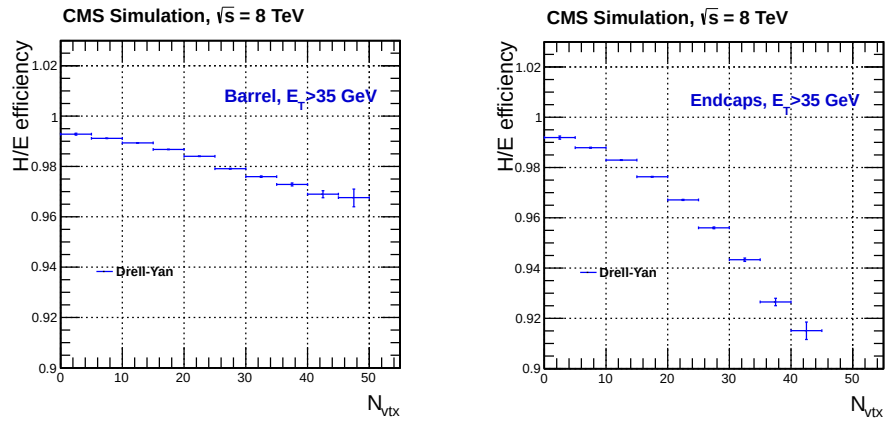
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(a)

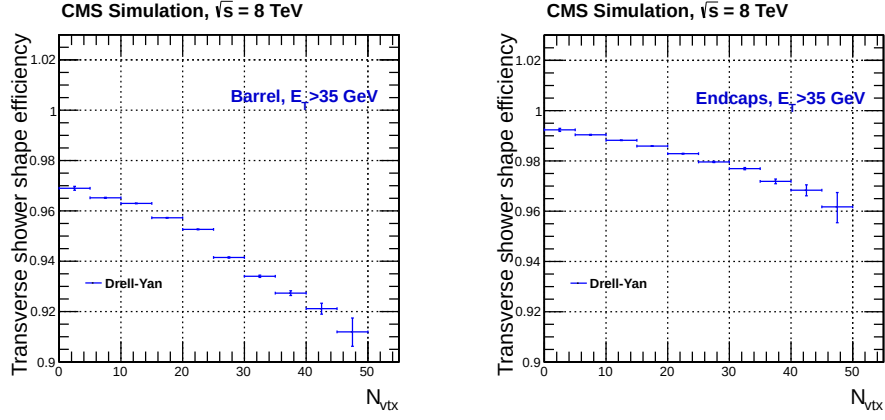


(b)

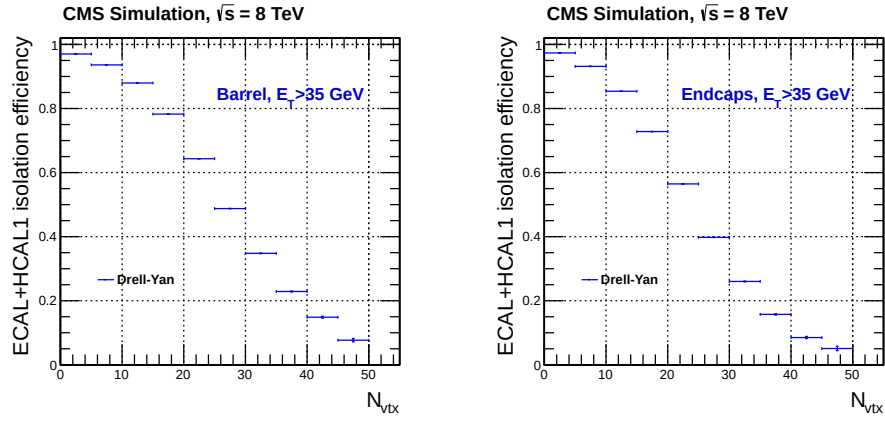


(c)

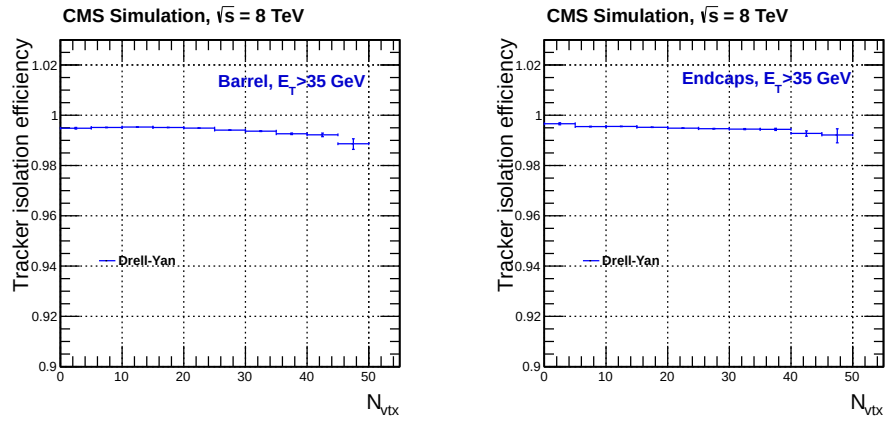
Figure 4.24: Efficiency of some of the cuts forming the high E_T electron selection as fonction of the number of vertices: $\Delta\eta_{in}$ (a), $\Delta\phi_{in}$ (b), H/E (c), in case of barrel (left) and endcaps (right) electrons with $E_T > 35$ GeV.



(a)



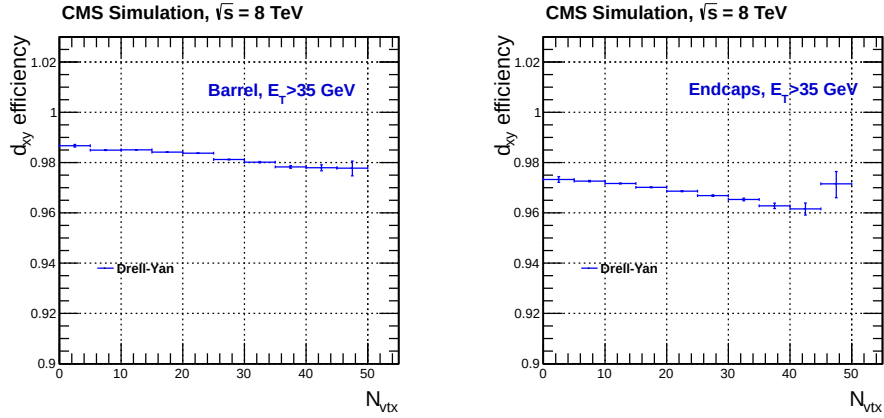
(b)



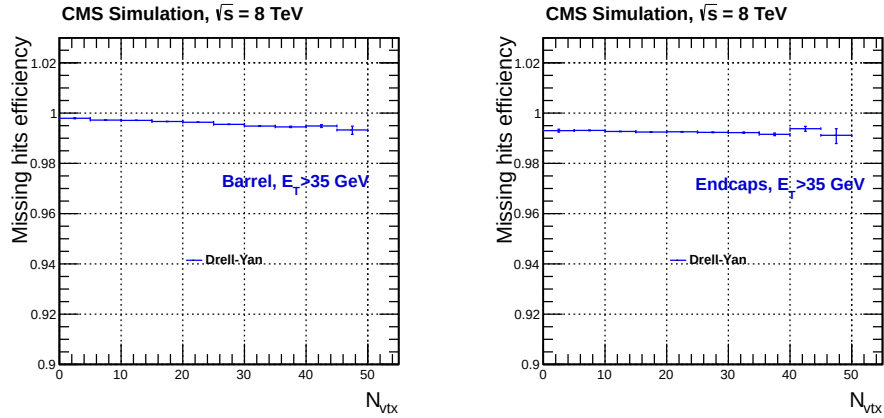
(c)

Figure 4.25: Efficiency of some of the cuts forming the high E_T electron selection as function of the number of vertices: the transverse shower shape cut (d), the ECAL+HCAL1 isolation (e), the tracker isolation (f), in case of barrel (left) and endcaps (right) electrons with $E_T > 35$ GeV.

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(a)



(b)

Figure 4.26: Efficiency of some of the cuts forming the high E_T electron selection as function of the number of vertices: d_{xy} (a), the number of missing hits (b), in case of barrel (left) and endcaps (right) electrons with $E_T > 35$ GeV.

of the mean isolation and of ρ as a function of the number of primary vertices:

$$E_A = \frac{\frac{d}{dN_{vtx}}(Mean\ Isolation)}{\frac{d}{dN_{vtx}}\rho} \quad (4.8)$$

This allows one to compute η dependent effective areas. Parametrizing the pile-up contamination using ρ rather than directly the number of vertices has two advantages. First, it gives a more accurate description of the activity due to pile-up events. Secondly it reduces the impact of a possible discrepancy in the pile-up activity between data and simulations. Although initially developed for the isolation, it is straightforward to extend the correction for any variable with dimensions of energy V :

$$V_{corr} = V_{raw} - \rho \times E_A \quad (4.9)$$

Corrections were therefore derived for the following variables:

- ECAL+HCAL1 Isolation
- $H/E \times E_T$
- The transverse shower shape variables in the barrel : $E_{5 \times 5}$, $E_{2 \times 5}$ and $E_{1 \times 5}$

For each cut under study an effective area needs to be calculated. This is performed using GSF electron candidates with $E_T > 35$ GeV coming from simulated events generated according to the Drell-Yan process, and matching a generated electron. The distribution of a variable including pile-up energy deposits sees its mean but also its width rising with the number of vertices and it is observed that the latter grows faster. As a consequence, the effective area depends on the targeted efficiency: for a high efficiency selection, the slope of the line connecting the values below which 95% of the distribution lies is therefore more appropriate than the slope of the distribution mean, and will in general give rise to higher E_A values.

4.4.3 ECAL+HCAL1 isolation correction

Figure 4.27 shows the mean values of the ECAL+HCAL1 and tracker isolation variables and ρ as a function of the number of primary vertices in the event separately for barrel and endcaps electrons. As expected all the curves have a linear behaviour with respect to the number of vertices: each extra vertex adds in average the same amount of energy to the isolation and the energy density. The tracker isolation actually shows a negligible dependency since only tracks compatible with the primary vertex of the electron are considered.

Figure 4.28 is a scatter plot of the ECAL+HCAL1 isolation as a function of the number of vertices. Each slice for a given number of vertices is normalized to 1. The blue curve connects the mean value of the isolation whereas the red curve shows the value below which exactly 95% of the events are found. This second curve is also close to a straight line and is rising steeper, meaning that the width of the distribution grows faster than the mean.

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The effective area, computed as the ratio of the slopes of a first order polynomial of the 95% efficient isolation cut value (red dots in figure 4.28) and ρ (black dots figure 4.27), is shown in figure 4.29 as a function of η in red for the ECAL and HCAL1 isolation separately and together. The same variable, computed this time as the ratio of the slopes of a first order polynomial of the mean isolation value (blue dots in figure 4.28) and ρ (black dots figure 4.27) as a function of the number of vertices, is drawn in black.

In the barrel, the ECAL+HCAL1 effective area is roughly flat for $|\eta| < 1$. The drop in the ECAL isolation observed at $|\eta| \approx 1.5$ is thought to be due to the fact that the isolation cone reaches the transition region between the ECAL barrel and endcap where energy leakage can happen. The increase of the HCAL1 effective area with $|\eta|$ can be explained by the fact that for a same p_t , particles in that region have more energy and will therefore penetrate further and are more likely to pass the noise threshold and be counted in the isolation sum.

As the η dependency is hard to model and rather small, it was decided to take a constant value. The original results, obtained from older simulations at $\sqrt{s} = 7$ TeV and shown in appendix C led to a similar shape with slightly higher values, resulting in an effective area of 0.28. While the effective area for the sum of the ECAL and HCAL1 isolations are roughly flat in η , it is not true for each calorimeter separately.

With this correction applied the associated cut in the high E_T electron selection becomes:

$$\begin{aligned} \text{Isol ECAL} + \text{HCAL1} &< 2 + 0.03 \times E_T + 0.28 \times \rho && (\text{Barrel}) \\ \text{Isol ECAL} + \text{HCAL1} &< 2.5 + 0.28 \times \rho && (\text{Endcaps, } E_T < 50 \text{ GeV}) \\ \text{Isol ECAL} + \text{HCAL1} &< 1 + 0.03 \times E_T + 0.28 \times \rho && (\text{Endcaps, } E_T > 50 \text{ GeV}) \end{aligned}$$

4.4.4 H/E correction

For GSF electron candidates with $E_T > 35$ GeV, the supercluster energy is largely dominated by the electron energy and the contribution of pile-up (a few GeV at most) has a very small effect on the H/E ratio. The H/E ratio is on the other hand very sensitive to the pile-up energy deposits in the HCAL. The same procedure than in section 4.4.3 was then applied to estimate the effective area corresponding to $H_T = H/E \times E_T$ (since ρ is a measurement of the average transverse energy density). The results are presented in figures 4.30, 4.31 and 4.32. A strong dependency with η is observed. The following function was found to fit fairly well this behaviour:

$$E_A(\eta) = A \times \cosh(B\eta) + C$$

It is drawn on top of figure 4.32, with the fitted values for A , B and C . The corresponding values obtained from the $\sqrt{s} = 7$ TeV simulations are 0.13, 0.41 and -0.085 (*cf* Appendix C). With this correction the H/E cut becomes:

$$H/E < 0.05 + (0.13 \times \cosh(0.41\eta) - 0.085) \times \rho/E_T$$

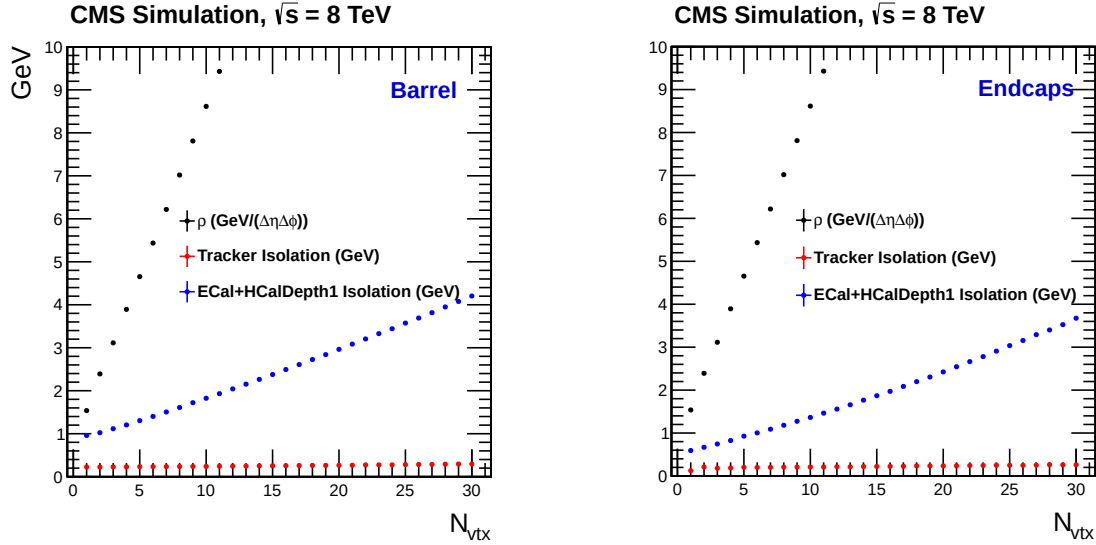


Figure 4.27: The mean value of the tracker isolation (red), ECAL+HCAL1 isolation (blue) as a function of the number of primary vertices in the event for electrons with $E_T > 35$ GeV. Also shown is the mean ρ value for such events.

4.4.5 Transverse shower shape cut correction

As a third correction, the subtraction of the pile-up contribution to the transverse shower shape variables in the barrel is studied. A correction in the endcaps is not investigated as the variable ($\sigma_{i\eta i\eta}$) is less affected by pile-up noise and more complicated to correct. The three variables to correct are $E_{5\times 5}$, $E_{2\times 5}$ and $E_{1\times 5}$. However the large electron energy (~ 50 GeV) makes very difficult the measurement of its pile-up dependency (of the order of 5 MeV per extra vertex). To get rid of most of the electron energy the effective area of the variable $E_{5\times 5,T} - E_{2\times 5,T}$ was rather computed. Since this effective area is calculated on a 2 domino, a geometrical scale factor of 5/3, 2/3 and 1/3 is then applied to correct $E_{5\times 5,T}$, $E_{2\times 5,T}$ and $E_{1\times 5,T}$ respectively.

The results for $E_{5\times 5,T} - E_{2\times 5,T}$ are shown on figures 4.33 and 4.34. Although the error bars are quite large, a positive value is measured that is roughly η independent. A flat line was fitted and gave an estimate of this effective area to be (0.0055 for $\sqrt{s} = 7$ TeV measurement). With this correction the associated cut in the selection then becomes:

$$E_{1\times 5,T} + 0.006 \times \rho > 0.83 \times E_{5\times 5,T} \parallel E_{2\times 5,T} + 0.005 \times \rho > 0.94 \times E_{5\times 5,T}$$

4.4.6 Electron selection efficiencies after pile-up correction

Figure 4.35 presents the efficiency as a function of the number of primary vertices in the event for each of the three corrected cuts for GSF electron candidates coming from

4. RECONSTRUCTION AND SELECTION OF HIGH ENERGY ELECTRONS WITH THE CMS DETECTOR

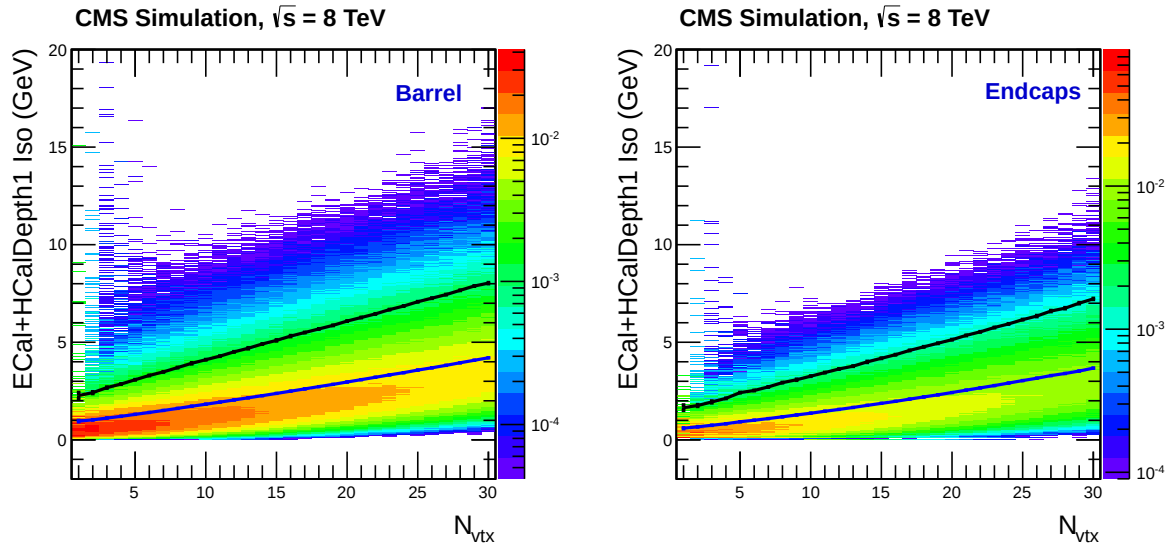


Figure 4.28: The distribution of the ECAL+HCAL1 Isolation versus the number of primary vertices in the event for electrons with $E_T > 35$ GeV. Each slice in constant number of vertices is normalized to 1. For a given number of vertices, the blue dots indicate the mean value of the projected distribution while the black dots indicate the value below which 95% of the events are contained.

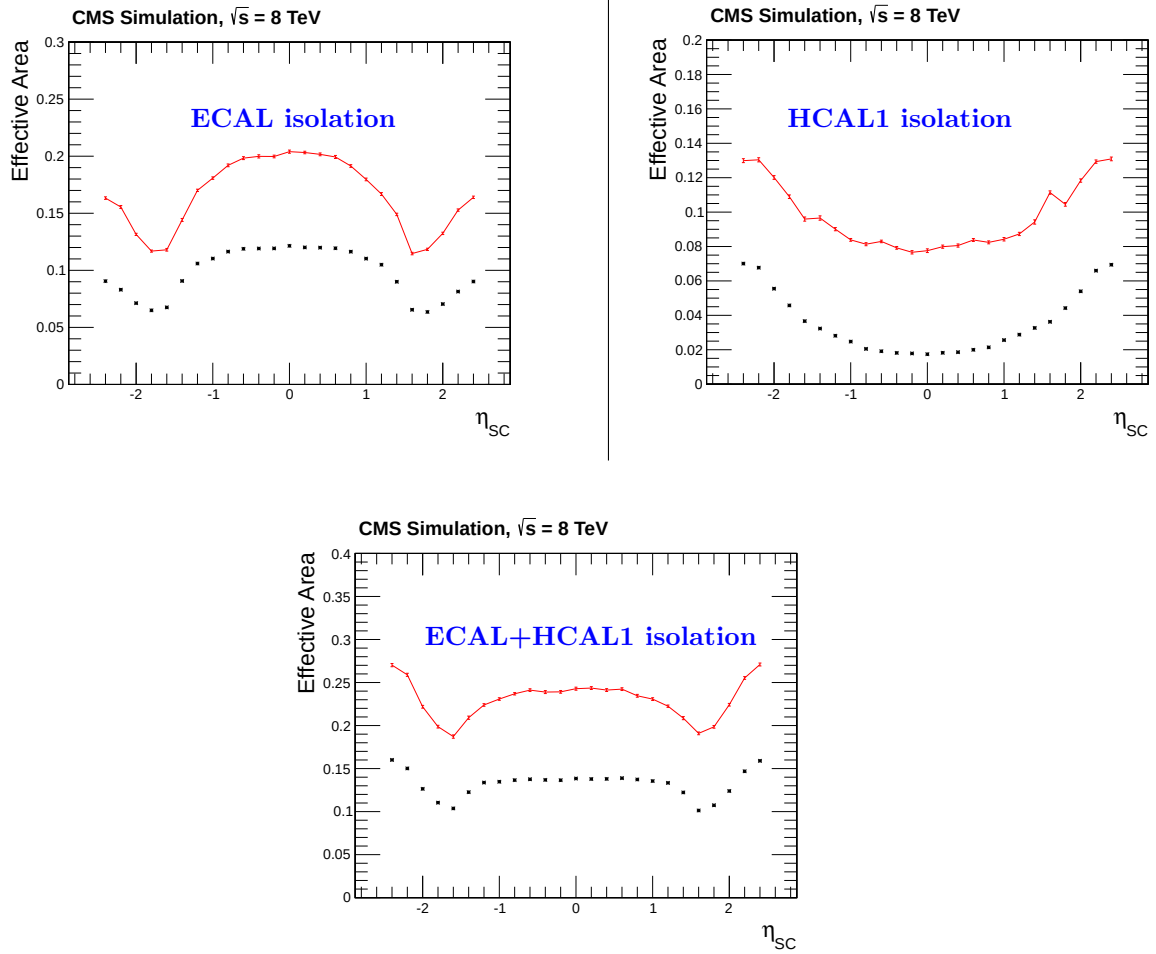


Figure 4.29: The effective areas for the ECAL and HCAL1 isolation variables separately (top) and together (bottom) as a function of the supercluster pseudorapidity obtained using the slopes of the mean isolation value (black) or the 95% efficient cut curve (red).

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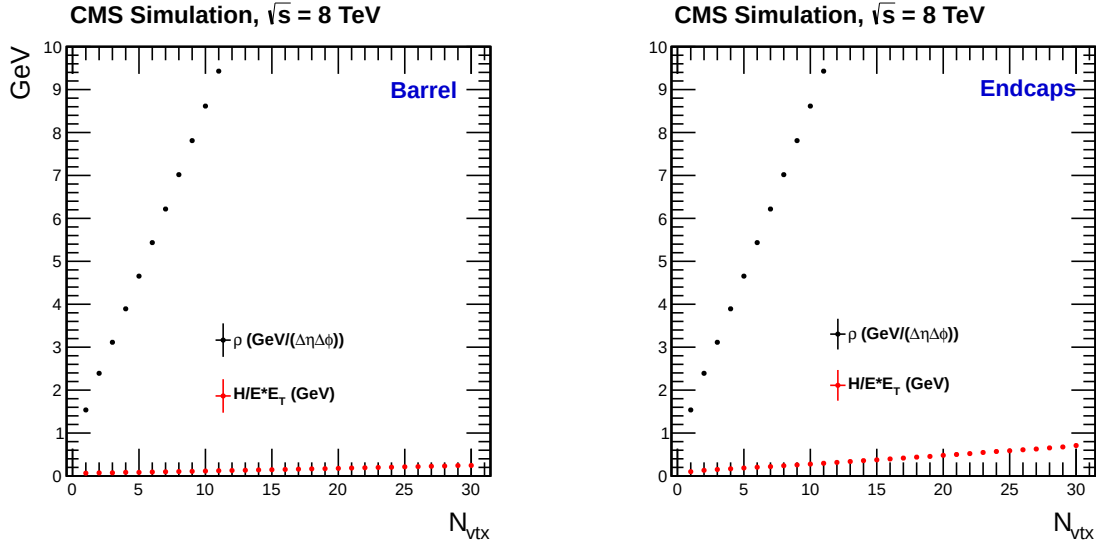


Figure 4.30: The mean value of the $H_T = H/E \times E_T$ (red) as a function of the number of primary vertices in the event for electrons with $E_T > 35$ GeV. Also shown is the mean ρ value for such events.

Drell-Yan events. The probability for fakes from QCD to pass the selection as a function of the number of reconstructed vertices was also studied and increases by around 50% between 1 and 50 vertices for $E_T = 35$ GeV. The dependency however disappears for E_T above 100 GeV. A small overcorrection is observed for the ECAL+HCAL1 isolation and H/E due to the fact that the effective areas calculated with simulations at $\sqrt{s} = 7$ TeV overestimate the ones obtained at $\sqrt{s} = 8$ TeV.

Only the correction on the ECAL+HCAL1 isolation was ultimately used in the final selection. The reason is that the shower shape corrections affect mainly moderate E_T values and are negligible for $E_T > 100$ GeV. They might however be used in the future.

4.5 Summary

This chapter details how the high energy electrons that will be used in the analysis presented in this thesis are reconstructed, triggered and identified with the CMS detector. The reconstruction algorithm, common for any kind of electron produced during a collision, consists of the construction of a GSF electron candidate by matching an energy deposit (supercluster) in the ECAL with a track reconstructed in the tracker (table 4.1). Electrons triggers start with the association of a supercluster with two hits in the first layers of the tracker, on top of which extra conditions are applied. The trigger dedicated to this analysis is detailed in table 4.2. The main part of this chapter presented a serie of additional cuts, designed for high energy electrons, that are applied to the GSF electron

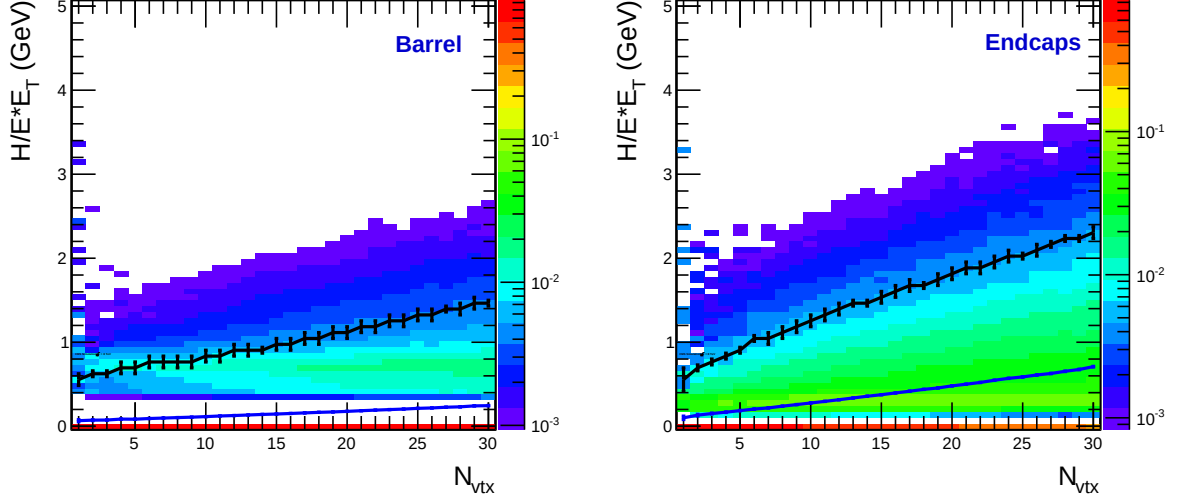


Figure 4.31: The distribution of the $H_T = H/E \times E_T$ versus the number of primary vertices in the event for electrons with $E_T > 35$ GeV. Each slice in constant number of vertices is normalized to 1. For a given number of vertices, the blue dots indicate the mean value of the projected distribution while the black dots indicate the value below which 95% of the events are contained.

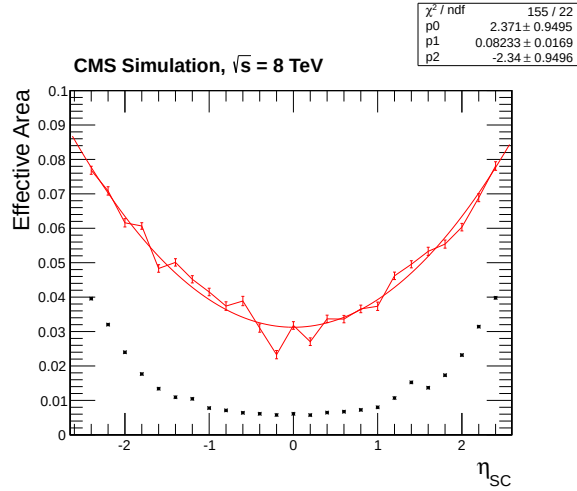


Figure 4.32: The Effective Area for the $H_T = H/E \times E_T$ variable as a function of the supercluster pseudorapidity obtained using the slopes of the mean $H_T = H/E \times E_T$ value (black) or the 95% efficient cut curve (red). The results of a fit to the parametric function $p0 \times \cosh(p1 \eta) + p2$ is also shown.

4. RECONSTRUCTION AND SELECTION OF HIGH ENERGY ELECTRONS WITH THE CMS DETECTOR

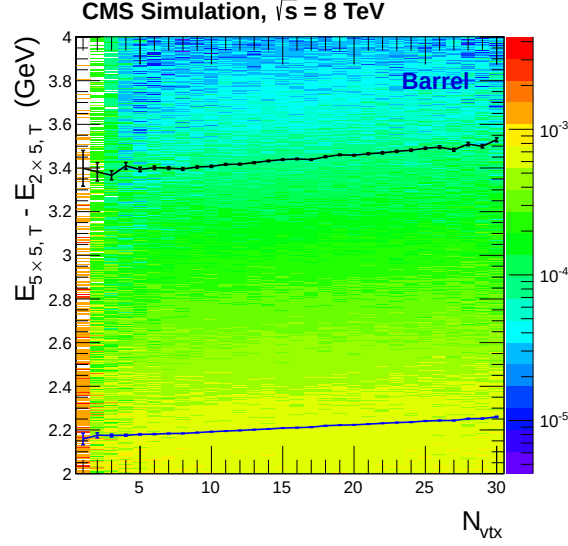


Figure 4.33: The distribution of the $E_{5 \times 5, T} - E_{2 \times 5, T}$ versus the number of primary vertices in the event for electrons with $E_T > 35$ GeV. Each slice in constant number of vertices is normalized to 1. For a given number of vertices, the blue dots indicate the mean value of the projected distribution while the black dots indicate the value below which 95% of the events are contained.

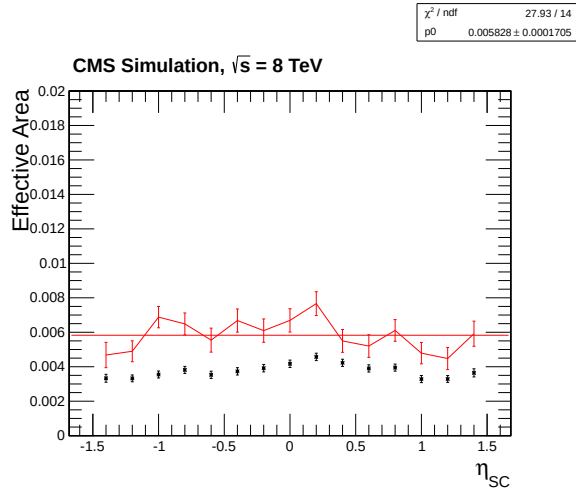
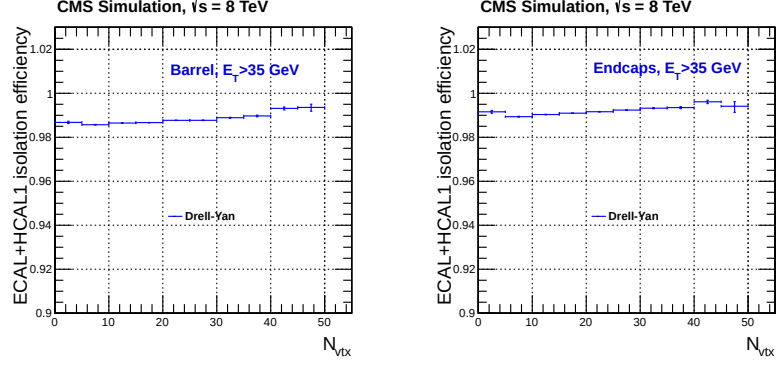
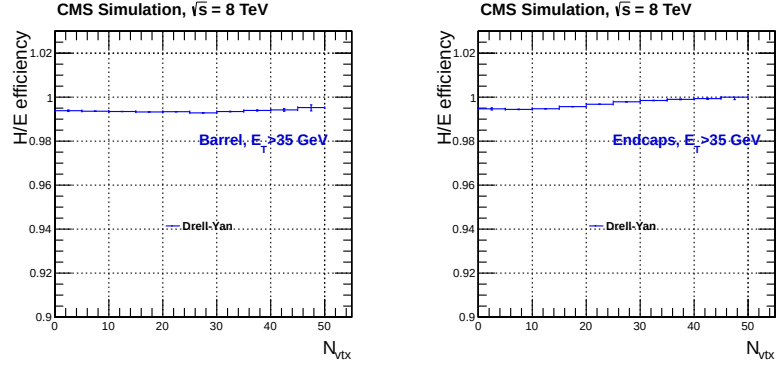


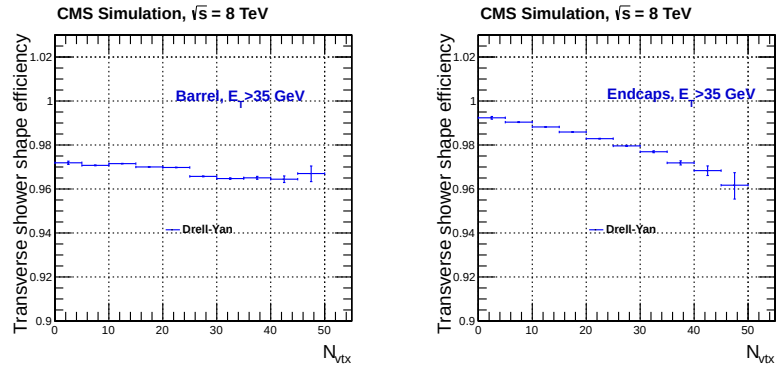
Figure 4.34: The Effective Area for the $E_{5 \times 5, T} - E_{2 \times 5, T}$ variable as a function of the supercluster pseudorapidity obtained using the slopes of the mean $E_{5 \times 5, T} - E_{2 \times 5, T}$ value (black) or the 95% efficient cut curve (red). The results of a fit to a straight line is also shown.



(a)



(b)



(c)

Figure 4.35: Efficiency of the ECAL+HCAL1 isolation (a), H/E (b) and the transverse shower shape (c) cuts applied on GSF electron candidates, after pile-up subtraction as a function of the number of primary vertices in case of barrel (left) and endcaps (right) electrons with $E_T > 35$ GeV.

4. RECONSTRUCTION AND SELECTION OF HIGH ENERGY ELECTRONS WITH THE CMS DETECTOR

Criteria	High E_T selection
SC-track matching	$ \Delta\eta_{in} < 0.005$ (0.007) $ \Delta\phi_{in} < 0.06$ (0.06)
Longitudinal shower shape	$H/E < 0.05$ (0.05)
Transverse shower shape	$E_{1\times 5}/E_{5\times 5} > 0.83$ OR $E_{2\times 5}/E_{5\times 5} > 0.94$ ($\sigma_{in\eta} < 0.03$)
Isolation	Tracker isolation < 5 GeV (5 GeV) AND ECAL+HCAL1 isolation - $0.28\times\rho < 2+0.03E_T$ ($< \max(2.5, 1+0.03E_T)$)
Vertex compatibility	$ d_{xy} < 0.02$ (0.05)
Photon rejection	# missing hits ≤ 1 (1)

Table 4.4: The variables used in the high E_T electron selection and the associated cut values for electrons emitted in the barrel (values in brackets indicate the corresponding cuts in the endcaps).

candidates to reduce the contamination from fakes. A special attention was paid to the impact of pile-up events on the selection efficiency and corrections factors were calculated. The final selection which includes a pile-up correction for the ECAL+HCAL1 variable is given in table 4.4. Its efficiency, for GSF electron candidates from Drell-Yan and QCD events is compared as a function of E_T in figure 4.36. Finally the complete efficiency including the reconstruction and selection steps is shown in figure 4.37 as a function of E_T for generated electrons from a Drell-Yan simulated sample emitted in the detector acceptance.

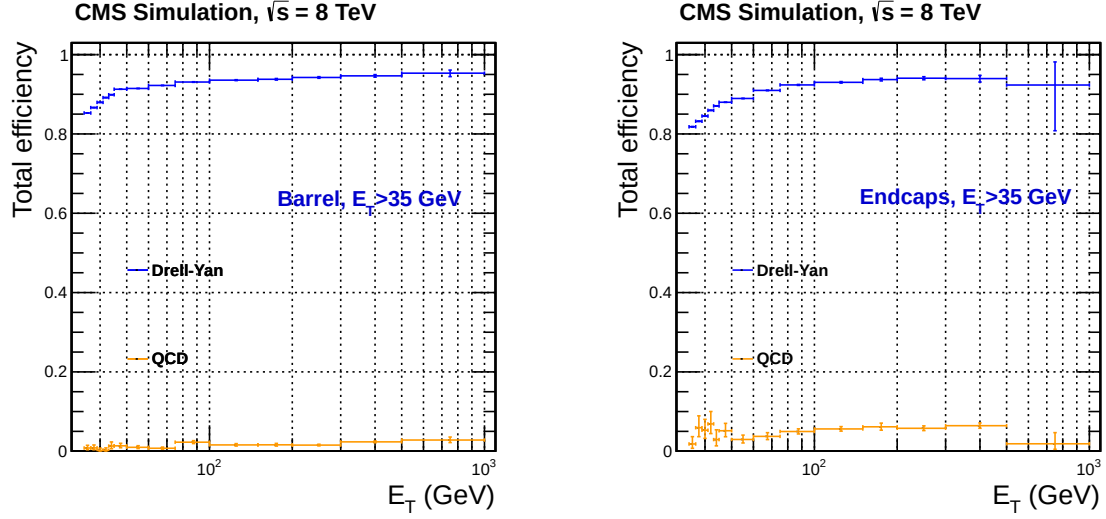


Figure 4.36: The total high E_T electron selection efficiency as a function of the GSF electron candidate E_T , for candidates from the Drell-Yan process (blue) and from QCD (orange).

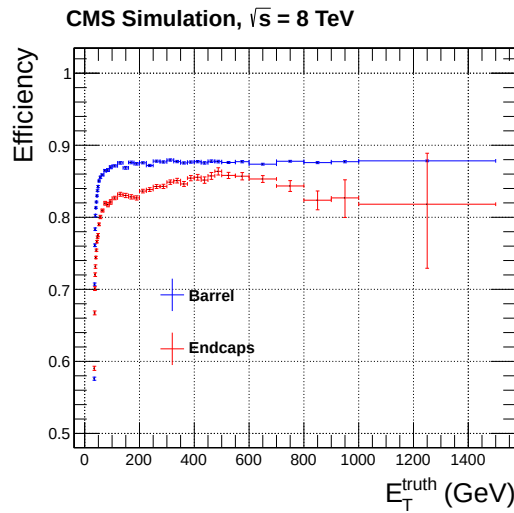


Figure 4.37: Full electron selection and reconstruction efficiency as a function of E_T for generated electrons from a Drell-Yan simulated sample emitted in the ECAL barrel (blue) or in the ECAL endcaps (red).

4. RECONSTRUCTION AND SELECTION OF HIGH ENERGY ELECTRONS WITH THE CMS DETECTOR

Chapter 5

Measurement of the high E_T electron selection, reconstruction and trigger efficiencies in the 2012 data

The reconstruction of GSF electron candidates and the high E_T electron presented in the previous chapter were designed using simulated samples. Any discrepancy between those simulation and the data can however strongly affect any physical analysis. A 5% discrepancy in the electron selection efficiency would for instance bias the cross section measurement of a process involving two electrons by 10%. It is therefore crucial to validate the electron reconstruction and selection performance with data. The search for heavy dielectron resonances performed in this thesis involves electrons with E_T up to 1 TeV. In that region, the lack of statistics in the data forbids to make a measurement. Rather, the efficiency and its uncertainty must be compared between data and simulation from medium values (35 GeV) to values as high as possible (500 GeV) and then extrapolated in the signal region using the simulation. A good confidence in the measurement in the highest E_T bins is therefore of primary importance.

The final event selection that will be used in chapter 6 requires two GSF electron candidates passing the high E_T electron selection. The event must also pass the double electron trigger described in table 4.2. The total efficiency for a generated electron to be finally selected can then be factorized into the following steps:

$$\epsilon_{SC} \times \epsilon_{GSF} \times \epsilon_{trigger} \times \epsilon_{offline\ selection} \quad (5.1)$$

where:

- ϵ_{SC} is the probability for an electron to form a supercluster in the ECAL,
- ϵ_{GSF} is the probability to find a track associated to the supercluster and to form a GSF electron candidate,
- $\epsilon_{trigger}$ is the probability for the GSF electron candidate to be reconstructed as a trigger object,
- $\epsilon_{offline\ selection}$ is the probability for the GSF electron candidate reconstructed as a trigger object to pass the high E_T electron selection.

5. MEASUREMENT OF THE HIGH E_T ELECTRON SELECTION, RECONSTRUCTION AND TRIGGER EFFICIENCIES IN THE 2012 DATA

In the present chapter, which mainly focuses on the high E_T electron selection, techniques to measure $\epsilon_{offline\ selection}$ are presented and applied to the full dataset recorded by CMS in 2012 [114]. This dataset is the same than the one used for the search of heavy resonances (*cf.* table 2.3). The techniques are also applied to events simulated according to the Drell-Yan process and a correction factor (“scale factor”) defined as the ratio of the data over simulation efficiency is computed.

In order to acquire confidence in the results at high transverse energy several complementary methods are used for the high E_T electron selection efficiency measurement.

Results for the $\epsilon_{trigger}$ and ϵ_{GSF} are also given.

The chapter is structured as follows: section 5.1 explains the principle of the so-called Tag and Probe method, section 5.2 presents the high E_T electron selection measurements obtained using dielectron pairs at the Z pole and section 5.3 is devoted to the same measurement using dielectron pairs in the high mass tail of the Drell-Yan process. Finally, results for the reconstruction efficiency and for the trigger efficiency are given in section 5.4

5.1 Description of the method and common event selection

The Tag and Probe method [115, 116] is designed to measure the efficiency of a given selection. It can be applied to various kinds of pair produced objects. The principle of the method is to require strong criteria on one of the two objects (referred to as the “tag”) and to measure the condition (here a cut) efficiency on the other one (the “probe”). The presence of a well reconstructed object in the event allows one to build a high purity sample where the contribution from fake probes is reduced. At the LHC, this technique is mainly applied to measure efficiencies associated to leptons with leptonic decays of J/ψ or Υ or with Drell-Yan (DY) events at the Z resonance. In this chapter, a data sample made of Drell-Yan events is built.

Assuming that there is no correlation between the two objects and the sample is pure, the efficiency is given by :

$$\epsilon = \frac{N_{TP}}{N_{TP} + N_{Tp}} \quad (5.2)$$

where N_{TP} (N_{Tp}) are the number of pairs where the probe passes (respectively fails) the cut under study. If several tag and probe pairs are found in an event, all of them are used in the efficiency measurement.

In sections 5.2 and 5.3, the tag and probe pairs consist of a GSF electron candidate with $E_T > 35$ GeV and passing the high E_T electron selection (the tag), and a GSF electron candidate with $E_T > 35$ GeV (the probe). The tag is in addition required to be reconstructed in the barrel where the fraction of fake electrons is much lower than in the endcaps. The event is also required to pass the trigger described in table 4.2 which requires two electrons with $E_T > 33$ GeV at the HLT and is based on a single e/γ candidate with $E_T > 22$ GeV at the L1. To avoid any bias in the measurement, the tag is finally matched with the L1 candidate: the L1 object and the probe must satisfy $\Delta R <$

0.5¹.

Before comparing the measurement made on electrons in the data with those in simulated Drell-Yan events, one must first subtract in the data the contribution from other processes which could bias the efficiency measurement.

Those backgrounds are of various types. First, there exists other processes with two electrons in the final state, such as the production of a $t\bar{t}$ pair, where both top quarks decay leptonically or the production of two heavy bosons (W or Z) with leptonic decays. Secondly, a sizeable fraction of the events selected in the data consists of pairs in which the probe is not an electron but a fake due to a misidentified jet. As discussed in section 4.3, fake GSF electron candidates are frequent in a hadron collider. Such pairs arise from several processes: W+jets events (where the tag is an electron coming from the W decay and the probe is a jet), γ +jets events (where the tag is usually a GSF electron candidate associated to a converted photon and the probe is a jet) and multijets events (where both the tag and the probe are jets). They will affect the efficiency measurement since the probability for a jet to pass the high E_T electron selection is much smaller than for an electron. All those backgrounds need to be properly estimated and subtracted in the data.

5.2 Selection efficiency measurement using events at the Z pole

In this section, the tag and probe pairs are required to have an invariant mass in the range 60 to 120 GeV where a large statistics can be collected. At leading order, such events are described by the Feynman diagram of figure 1.2 and the two electrons satisfy $E_T \leq M_Z/2$ GeV. Probing higher electron E_T values can however only be reached if the Z boson is boosted in the transverse plane. This happens for events described by higher order Feynman diagrams such as those shown in figures 3.4 and 3.5, which include at least one hard jet to balance the Z boson transverse momentum. Such events are therefore rarer.

The residual background in the data sample can be estimated either using simulation or taking advantage of the difference in shape between the invariant mass distribution of Drell-Yan events and the non resonant background. Both approaches are studied in the following sections.

5.2.1 Background subtraction from simulation

The invariant mass spectrum of all the tag and probe pairs is shown in figure 5.1 (pairs with two tags are counted twice) where the data (black dots) are compared to the expectations (stacked histogram) which are split into three categories: Drell-Yan events (red), other processes with two prompt electrons (green), processes where at least one electron is a fake (yellow). The expectations are obtained from simulation for all the processes except for the multijet contributions which is estimated using the “fake rate” technique detailed in section 6.3.3: the fake rate, denoted ζ , and defined as the probability, for a GSF electron candidate associated to a jet, to pass the high E_T electron selection, is measured in the

¹ The reason for such a loose ΔR cut value comes from the fact that L1 objects can only have fixed and coarsely spaced η and ϕ values.

5. MEASUREMENT OF THE HIGH E_T ELECTRON SELECTION, RECONSTRUCTION AND TRIGGER EFFICIENCIES IN THE 2012 DATA

data as a function of the GSF electron candidate η and E_T . Since the tag is simply a GSF electron candidate passing the electron selection and emitted in the barrel, the contribution of multijets to the tag and probe pairs is therefore obtained from pairs of GSF electron candidates labeled 1 and 2 as:

$$N_{TP,multijets} = N_{pp} \cdot \frac{\zeta(\eta_1, E_{T1})}{(1 - \zeta(\eta_1, E_{T1}))} \frac{\zeta(\eta_2, E_{T2})}{(1 - \zeta(\eta_2, E_{T2}))} (\theta_1 + \theta_2) \quad (5.3a)$$

$$N_{Tp,multijets} = N_{pp} \cdot \frac{\zeta(\eta_1, E_{T1})}{(1 - \zeta(\eta_1, E_{T1}))} \theta_1 + N_{pp} \frac{\zeta(\eta_2, E_{T2})}{(1 - \zeta(\eta_2, E_{T2}))} \theta_2 \quad (5.3b)$$

where N_{pp} is the number of pairs of GSF electron candidates in the data where both candidates fail to pass the electron selection, $\zeta(\eta_1, E_{T1})$ and $\zeta(\eta_2, E_{T2})$ are the fake rates associated to each of the candidate, and θ_i is 1 if the candidate i is emitted in the barrel and 0 otherwise. The two equations 5.3 are the sum of two terms where each GSF electron candidate plays successively the role of the tag or the probe. The denominators relate the total number of pairs to N_{pp} . The use of N_{pp} ensures to work with a sample completely dominated by fake electrons and where the contribution from Drell-Yan process is negligible.

The small amount of events observed in figure 5.1 for masses below 80 GeV is the consequence of a kinematical effect: this region is disfavoured because of the E_T threshold. The number of pairs in the data and the expected fractions of the different processes are shown, for different probe E_T ranges, in tables 5.1 and 5.2 separately for probes reconstructed in the barrel or in the endcaps. The main backgrounds that will affect the measurement come from events where the probe is actually a jet while the tag is a true electron (W+jets) or not (γ +jets, multijets). The $t\bar{t}$ and diboson background also contribute a little but both the tag and the probes are in this case electrons and they are not expected to strongly affect the efficiency, although the possible presence of jets in the event could reduce the selection efficiency if an electron is emitted close to one of the jets. As most of the electrons from the Drell-Yan process have E_T values close to half the Z mass, the background fraction increases for high E_T values and reaches a maximum of around 10% for in the range 70-100 GeV due to the simultaneous conditions on the electrons E_T and the dielectron invariant mass. The total number of events is rapidly decreasing when probing high E_T since this implies working with Z bosons strongly boosted in the transverse plan. Finally, in the endcaps, the phase space is further reduced at very high E_T because of the condition that the tag electron lies in the barrel (*cf.* last column of tables 5.1 and 5.2).

The selection efficiency is computed as a function of various variables by counting the number of tag and probe pairs where the probe passes the high E_T electron selection over the total number of tag and probe pairs and after subtracting the expected contribution from background to the numerator and the denominator. The scale factor is defined as the ratio of the efficiencies measured in the data after background subtraction and measured in a simulated sample of Drell-Yan events. To take into account the uncertainty on the background contribution, a systematic error, computed by varying the background contribution by 40%, is assigned to the scale factor. This reflects the theoretical uncertainty on the cross sections of the various processes and in particular on the Drell-Yan process

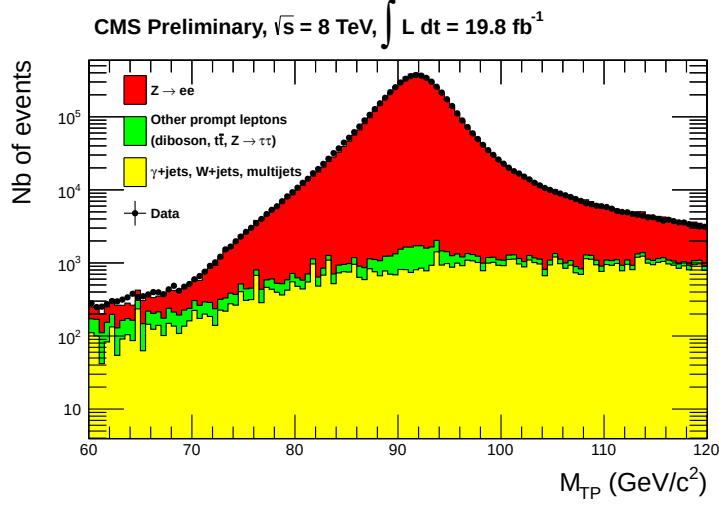


Figure 5.1: Invariant mass of tag and probe pairs with the selection at Z peak. Events with two tags are counted twice. The complete 2012 dataset corresponding to 19.8 fb^{-1} is used. The number of events in the simulation is normalized to the number of events in the data.

Barrel					
E_T range (GeV)	35-50	50-70	70-100	100-300	> 300
Number of events (data)	3947550	670667	55225	21503	341
Drell-Yan	98.4%	96.6%	89.6%	92.0%	93.4%
Other prompt leptons	0.3%	1.0%	4.4%	4.9%	6.0%
Jets	1.2%	2.4%	6.0%	3.1%	0.6%

Table 5.1: Number of events in data and expected fractions of the different processes for tag and probe pairs with $60 < M_{TP} < 120 \text{ GeV}$ for probes reconstructed in the barrel for various probe E_T ranges.

Endcaps					
E_T range (GeV)	35-50	50-70	70-100	100-300	> 300
Number of events (data)	815698	83322	8425	2957	9
Drell-Yan	98.5%	95.0%	91.8%	95.4%	95.0%
Other prompt leptons	0.3%	1.1%	3.7%	3.7%	3.8%
Jets	1.2%	3.9%	4.5%	0.9%	1.1%

Table 5.2: Number of events in data and expected fractions of the different processes for tag and probe pairs with $60 < M_{TP} < 120 \text{ GeV}$ for probes reconstructed in the endcaps for various probe E_T ranges.

5. MEASUREMENT OF THE HIGH E_T ELECTRON SELECTION, RECONSTRUCTION AND TRIGGER EFFICIENCIES IN THE 2012 DATA

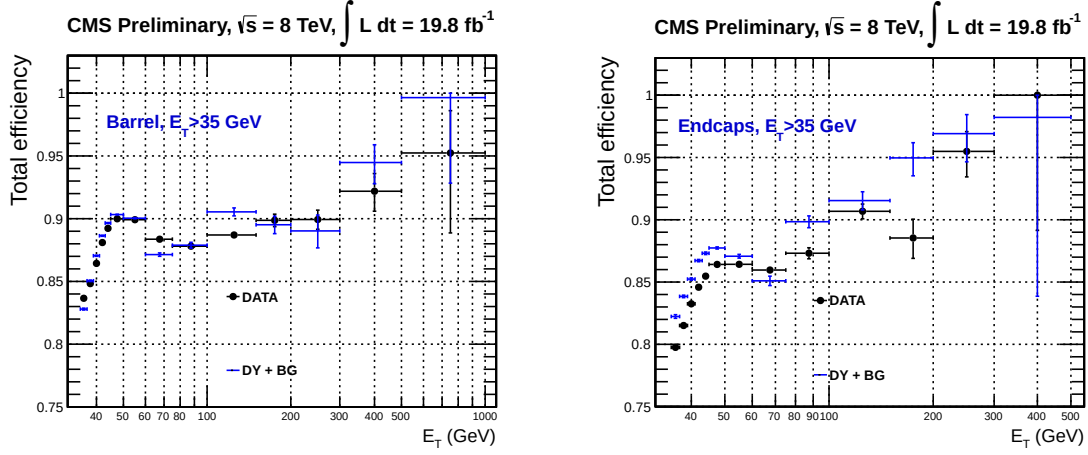


Figure 5.2: High E_T electron selection efficiency as a function of E_T for probe electrons with $E_T > 35$ GeV measured in a simulated Drell-Yan + background sample (blue) and in the data (black dots) before background subtraction. Efficiencies for electrons reconstructed in the barrel (left) and endcaps (right) are shown separately. Only statistical errors are drawn for the efficiencies.

in the different E_T bins considered [117, 118, 119] and the experimental precision on the fake rate measurement (as detailed in section 6.3.3).

The efficiencies are presented in figure 5.2 as a function of the probe E_T , for the data before background subtraction (black triangles) and a simulated sample containing Drell-Yan and background events (blue). One observes that the efficiency first increases up to 50 GeV, then drops to reach a local minimum around 70 GeV and finally increases again. This behaviour is due to the variation of the background contamination in the different E_T bins (*cf.* tables 5.1 and 5.2).

Figure 5.3 shows the same efficiency after background subtraction in the data and compares it to the expectation obtained from Drell-Yan simulated events. As expected, the efficiency now increases with E_T on the whole range and reaches a plateau at 95% in the barrel. The scale factor is also drawn in figure 5.3. In the barrel a value close to 1 is measured while in the endcaps, it is rather 0.98. The source of the discrepancy in the endcaps was identified to come mainly from an imperfect description of the tracker material budget leading to a spread larger than expected of the $\Delta\eta_{in}$ variable in the data. This effect impacts especially electrons close to the ECAL barrel and endcaps transition region. The fact that the scale factor is roughly constant in the whole E_T range gives however confidence that the scale factor can be extrapolated to very high E_T .

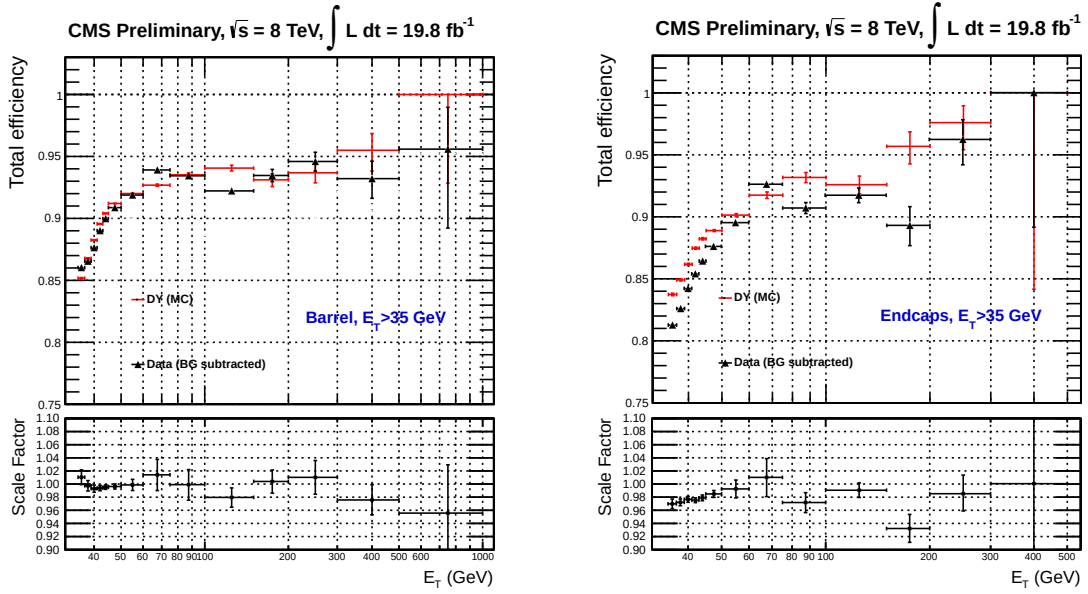


Figure 5.3: High E_T electron selection efficiency and scale factor as a function of E_T for probe electrons with $E_T > 35$ GeV measured in a simulated Drell-Yan only sample (red), and in the data after background subtraction (black triangles). Efficiencies for electrons reconstructed in the barrel (left) and endcaps (right) are shown separately. Only statistical errors are drawn for the efficiencies. Both statistical and systematic uncertainties are drawn for the scale factor.

5. MEASUREMENT OF THE HIGH E_T ELECTRON SELECTION, RECONSTRUCTION AND TRIGGER EFFICIENCIES IN THE 2012 DATA

	Barrel, $E_T \geq 35$ GeV	Endcaps, $E_T \geq 35$ GeV
Efficiency (Drell-Yan)	$0.897 \pm 0.000(stat.)$	$0.870 \pm 0.000(stat.)$
Efficiency (Drell-Yan + BG)	$0.884 \pm 0.000(stat.)$	$0.858 \pm 0.000(stat.)$
Efficiency (Data)	$0.882 \pm 0.000(stat.)$	$0.840 \pm 0.000(stat.)$
Efficiency (Data BG corr.)	$0.895 \pm 0.000(stat.)$	$0.853 \pm 0.000(stat.)$
Scale factor	$0.997 \pm 0.000(stat.) \pm 0.006(syst.)$	$0.979 \pm 0.000(stat.) \pm 0.006(syst.)$
	Barrel, $E_T \geq 100$ GeV	Endcaps, $E_T \geq 100$ GeV
Efficiency (Drell-Yan)	$0.939 \pm 0.002(stat.)$	$0.933 \pm 0.006(stat.)$
Efficiency (Drell-Yan + BG)	$0.904 \pm 0.002(stat.)$	$0.923 \pm 0.006(stat.)$
Efficiency (Data)	$0.891 \pm 0.002(stat.)$	$0.906 \pm 0.005(stat.)$
Efficiency (Data BG corr.)	$0.925 \pm 0.002(stat.)$	$0.915 \pm 0.005(stat.)$
Scale factor	$0.985 \pm 0.002(stat.) \pm 0.014(syst.)$	$0.981 \pm 0.006(stat.) \pm 0.004(syst.)$

Table 5.3: High E_T electron selection efficiency for electrons reconstructed in the barrel (left) and endcaps (right), obtained from the Tag and Probe method at the Z pole on a simulated Drell-Yan sample, on a simulated Drell-Yan and background sample, and on the 2012 data, before and after background subtraction. The last row gives the scale factor, taken as the ratio of the efficiency measured in data after background subtraction to the one measured in the simulated Drell-Yan sample. The upper (lower) part of the table shows the results for probe with $E_T > 35$ GeV (100 GeV). The systematic error on the scale factor is obtained by varying the expected background contribution by 40%.

The same efficiency and scale factor as in figure 5.3 are presented in figures 5.4 and 5.5 as a function of the probe η and ϕ and the number of reconstructed primary vertices in the event respectively.

The drop of efficiency close to the ECAL barrel-endcaps gap ($|\eta| \approx 1.5$) in the data is visible in figure 5.4. Small variations of the efficiency in ϕ , well reproduced by the simulation, can also be observed. As discussed in chapter 4, a small dependency of the efficiency versus the number of vertices is expected (figure 5.5). This effect is indeed observed in the data and well modeled by the simulation. It vanishes for electrons with $E_T > 100$ GeV as demonstrated in the bottom plot of figure 5.5.

Finally, the efficiencies and scale factor values for electrons with $E_T > 35$ GeV and $E_T > 100$ GeV are listed in table 5.3.

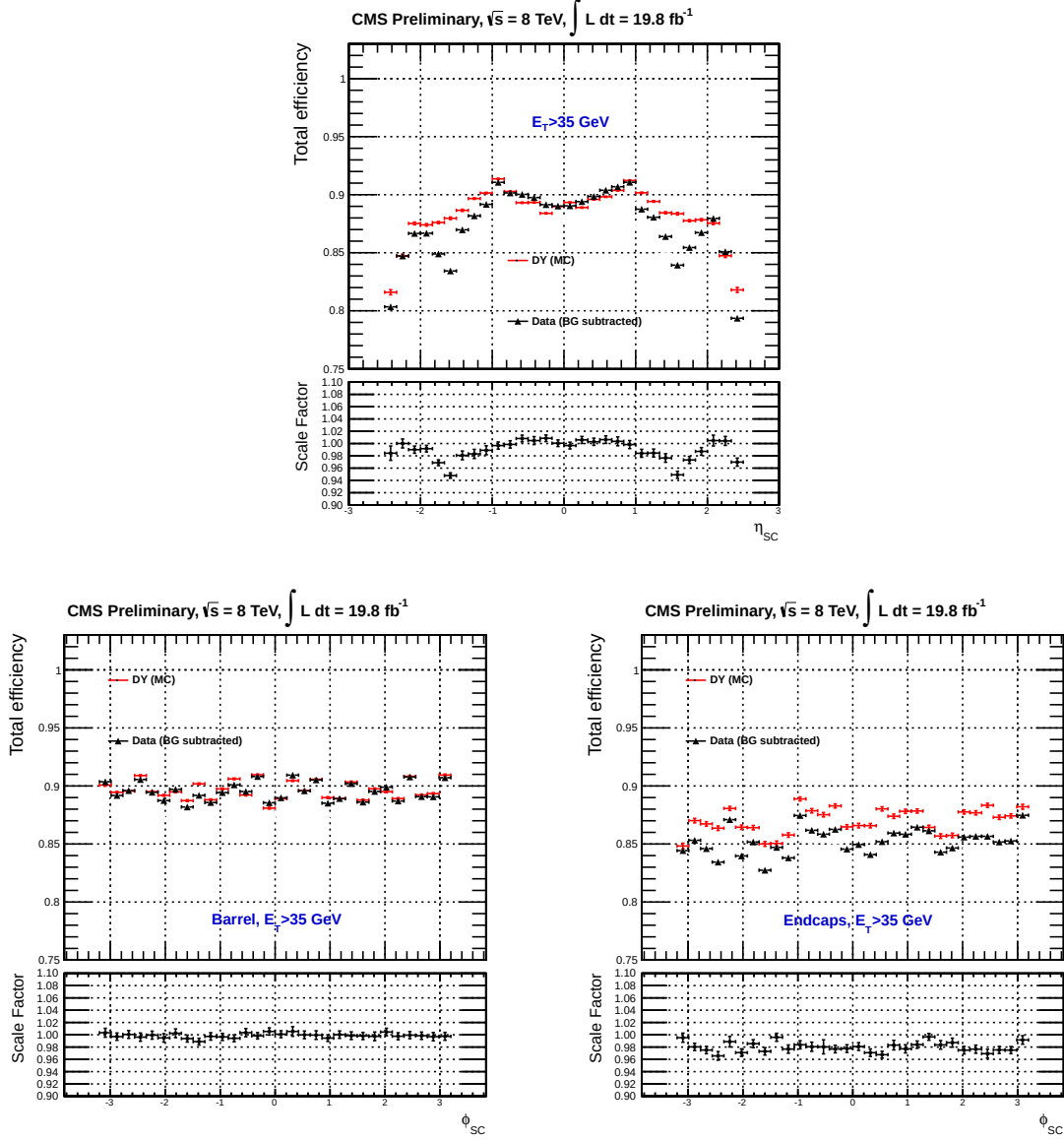


Figure 5.4: High E_T electron selection efficiency and scale factor as a function of η (top) and ϕ (bottom) for probe electrons with $E_T > 35$ GeV. Efficiencies as a function of ϕ are shown separately for electrons reconstructed in the barrel (left) and endcaps (right). The efficiency is shown for a simulated Drell-Yan sample (red) and for the data after background subtraction (black). Only statistical errors are drawn for the efficiencies. Both statistical and systematic uncertainties are drawn for the scale factor.

5. MEASUREMENT OF THE HIGH E_T ELECTRON SELECTION, RECONSTRUCTION AND TRIGGER EFFICIENCIES IN THE 2012 DATA

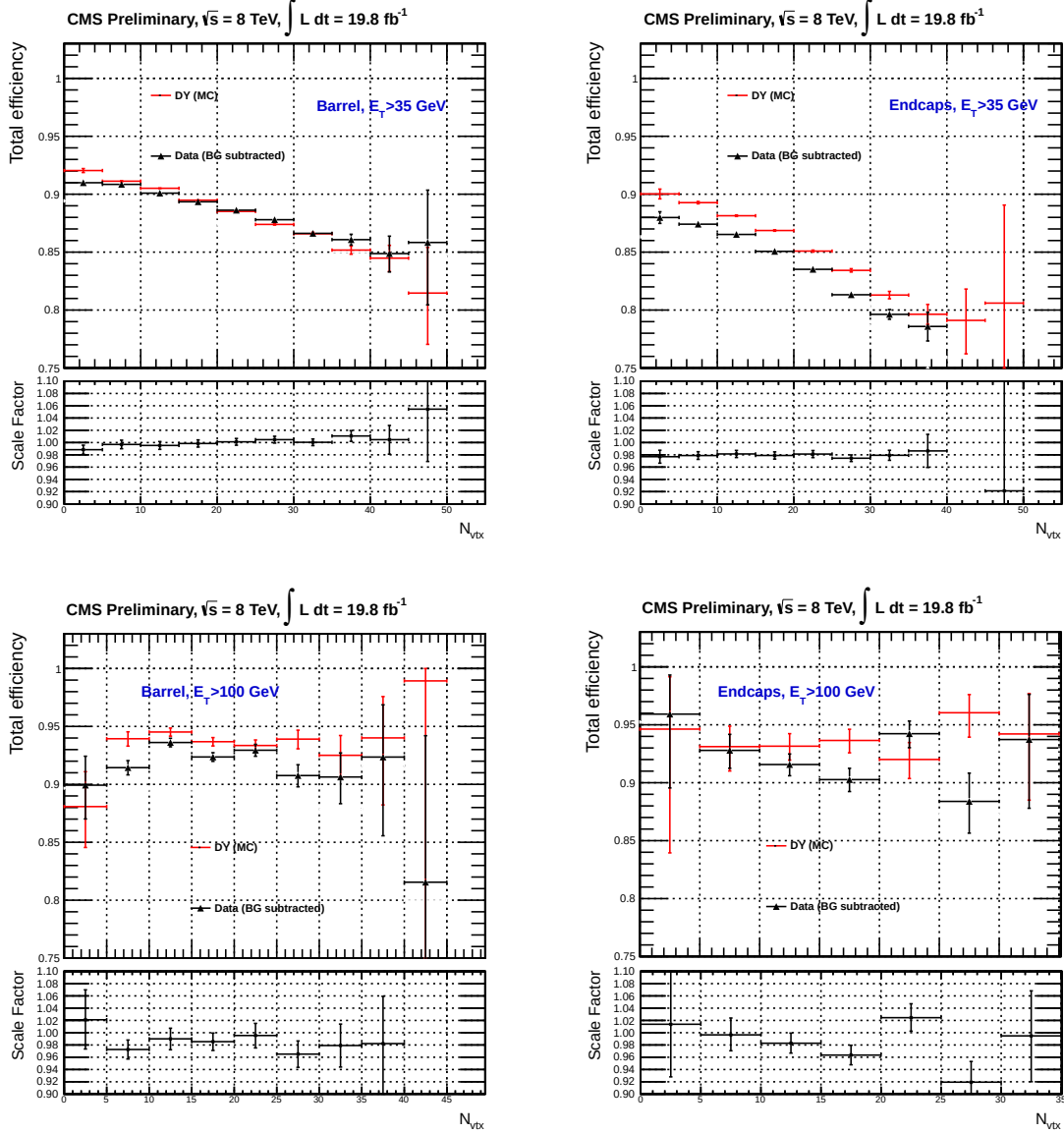


Figure 5.5: High E_T electron selection efficiency and scale factor as a function of the number of primary vertices in the event for probe electrons with $E_T > 35$ GeV (top) and $E_T > 100$ GeV (bottom). Efficiencies for electrons reconstructed in the barrel (left) and endcaps (right) are shown separately. The efficiency is shown for a simulated Drell-Yan sample (red) and for the data after background subtraction (black). Only statistical errors are drawn for the efficiencies. Both statistical and systematic uncertainties are drawn for the scale factor.

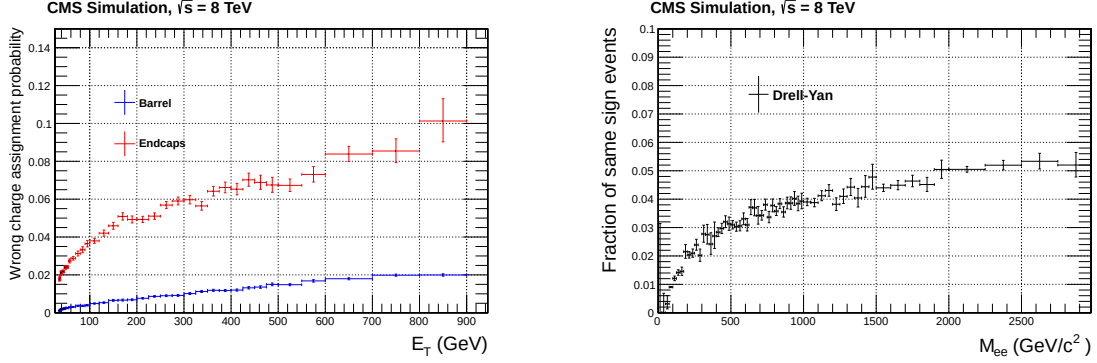


Figure 5.6: The wrong charge assignment probability for electrons from a Drell-Yan simulated sample as a function of their transverse energy (left), for GSF electron candidates reconstructed in the ECAL barrel (blue) or endcaps (red).

5.2.1.1 Wrong charge assignment

In addition to the high E_T electron selection, the probability of assigning the bad charge to a reconstructed GSF electron candidate was also studied.

For high E_T electrons, the track trajectory is close to a straight line in the transverse plane and a mismeasurement of the sign of its curvature can happen. This occurs in particular when a bremsstrahlung photon converts into a dielectron pair. One of the secondary electrons can leave hits in the tracker detector that are sometimes used in the reconstruction of the GSF track of the primary electron. If this secondary electron carries a charge opposite to the one of the primary electron, the wrong charge can be measured for the GSF electron candidate. The wrong charge assignment probability for candidates passing the E_T electron selection is shown as a function of the generated electron E_T from a Drell-Yan simulated sample separately for barrel and endcaps candidates in figure 5.6 (left) as well as the fraction of dielectron pairs reconstructed with a same charge as a function of the generated invariant mass of the pair (right). This measurement error can reach 10% in the endcaps: the size of the electron trajectory is shorter than in the barrel and the estimation of its curvature is therefore harder. For dielectron pairs with an invariant mass above 2 TeV, 5% of the pairs consist of two electrons reconstructed with the same charge.

These results on simulation were checked on the data using the tag and probe method. If one denotes $p_{T,WC}$ and $p_{P,WC}$ the charge misidentification probability for the tag and the probe electron respectively, the ratio of the number of pairs where the tag and the probe are reconstructed with a same charge, N_{SS} , over the number where they are reconstructed with an opposite charge, N_{OS} , is given by:

$$\frac{N_{SS}}{N_{OS}} = \frac{p_{T,WC} \times (1 - p_{P,WC}) + (1 - p_{T,WC}) \times p_{P,WC}}{(1 - p_{T,WC}) \times (1 - p_{P,WC}) + p_{T,WC} \times p_{P,WC}} \quad (5.4)$$

If $p_{T,WC} \ll p_{P,WC} \ll 1$, $p_{P,WC}$ can be approximated by $\frac{N_{SS}}{N_{OS}}$ and this probability can then be studied as a function of E_T , η , ... The assumption concerning the charge of the tag

5. MEASUREMENT OF THE HIGH E_T ELECTRON SELECTION, RECONSTRUCTION AND TRIGGER EFFICIENCIES IN THE 2012 DATA

can be satisfied by using the strong tag selection described later on in table 5.5. For such electrons with $E_T \approx 35$ GeV, the mischarge probability is expected to be around 0.05% against 0.2% for electrons passing the high E_T selection and emitted in the barrel. This is mainly due to the tight cuts on the $\Delta\phi_{in}$ and E/p which remove electrons emitting bremsstrahlung and which have more chance to be assigned the wrong charge.

Figures 5.7 and 5.8 show the fraction of tag and probe pairs where the two electrons carry an opposite charge in the Z peak window and pass the high E_T electron selection. The background for this selection is expected to be negligible and this allows one to directly compare the data to electrons from the Drell-Yan process. Figures 5.9 and 5.10 show the same variable for pairs where the probe is a simple GSF electron candidate. In this case, the background is not negligible anymore and those distributions can be used to check the background estimation. The good agreement between the data and the simulation in both sets of plots gives confidence that both the mischarge probability and the background contamination are well estimated by the simulation.

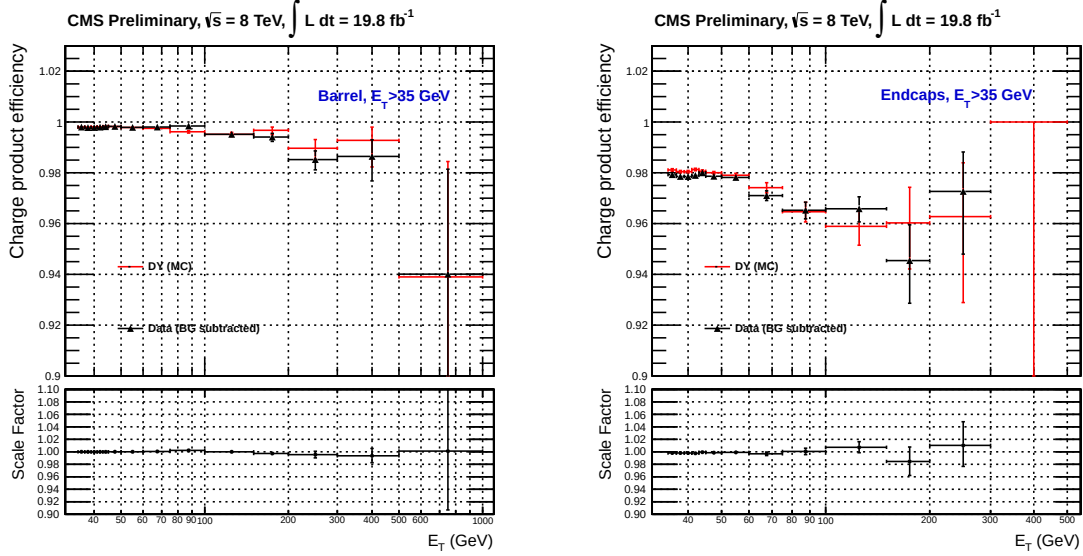


Figure 5.7: Fraction of opposite sign tag and probe pairs as a function of the probe E_T for tag and probe pairs where the probe passes the high E_T electron selection and the tag passes the selection described in table 5.5. This is the probability for the probe to be assigned the wrong charge if $p_{T,WC} \ll p_{P,WC}$. Results are shown separately for probes emitted in the barrel (left) or in the endcaps (right).

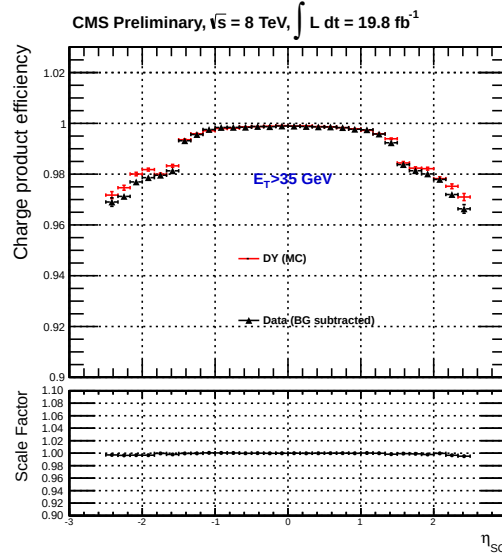


Figure 5.8: Fraction of opposite sign tag and probe pairs as a function of the probe η for tag and probe pairs where the probe passes the high E_T electron selection and the tag passes the selection described in table 5.5. This is the probability for the probe to be assigned the wrong charge if $p_{T,WC} \ll p_{P,WC}$.

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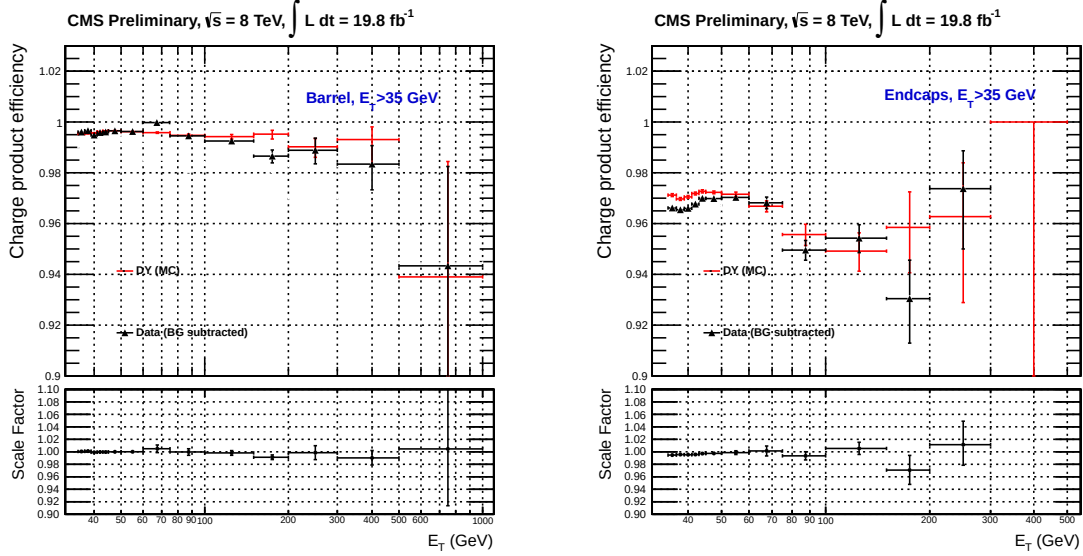


Figure 5.9: Fraction of opposite sign tag and probe pairs as a function of the probe E_T for tag and probe pairs where the probe is GSF electron candidate and the tag passes the selection described in table 5.5. This is the probability for the probe to be assigned the wrong charge if $p_{T,WC} \ll p_{P,WC}$. Results are shown separately for probes emitted in the barrel (left) or in the endcaps (right).

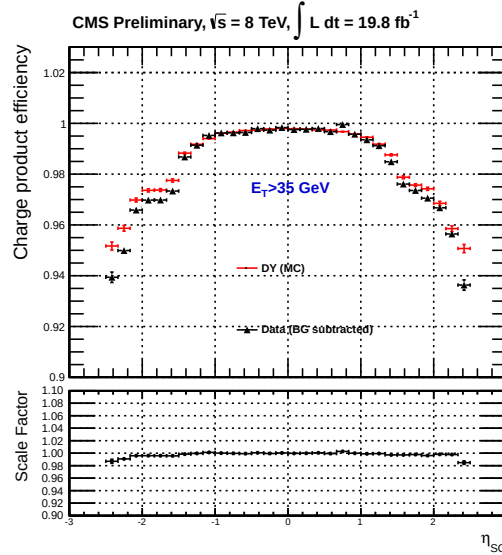


Figure 5.10: Fraction of opposite sign tag and probe pairs as a function of the probe η for tag and probe pairs where the probe is GSF electron candidate and the tag passes the selection described in table 5.5. This is the probability for the probe to be assigned the wrong charge if $p_{T,WC} \ll p_{P,WC}$.

5.2.2 Background subtraction with a fitting method at the Z peak

In order to cross check the results of section 5.2.1, the measurement of the high E_T electron selection efficiency was also performed with a background estimation directly from the data using a fitting method. The principle is to take advantage of the difference in shape between the tag and probe mass distributions for signal and background events in the Z peak region to estimate their respective yields. The tag and probe pairs are separated into two categories depending on whether the probe passes the cut under study (passing pairs) or not (failing pairs). A fit on the mass (M) distribution of the passing and failing pairs is simultaneously performed using a likelihood minimization, in which the total number of events is allowed to float. The probability density function (pdf) is binned and described for the passing pairs by:

$$N_{sig,pass} \times f_{sig,pass}(M, \vec{\alpha}_{sig,pass}) + N_{bkg,pass} \times f_{bkg,pass}(M, \vec{\alpha}_{bkg,pass}) \quad (5.5)$$

and for the failing pairs by:

$$N_{sig,fail} \times f_{sig,fail}(M, \vec{\alpha}_{sig,fail}) + N_{bkg,fail}(M) \times f_{bkg,fail}(M, \vec{\alpha}_{bkg,fail}) \quad (5.6)$$

where $N_{sig/bkg,pass/fail}$ are the estimated numbers of passing or failing pairs from Drell-Yan process (signal) and from all the other processes (background) and $f_{sig/bkg,pass/fail}$ the corresponding pdf described by a set of parameters $\vec{\alpha}_{sig/bkg,pass/fail}$. These definitions allow one to directly put the efficiency as a parameter in the fit using the relation:

$$\epsilon_{sig/bkg} = \frac{N_{sig/bkg,pass}}{N_{sig/bkg,pass} + N_{sig/bkg,fail}} \quad (5.7)$$

The complete set of parameters entering the fits are therefore $\epsilon_{sig/bkg}$, $N_{sig/bkg,TOT} = N_{sig/bkg,pass} + N_{sig/bkg,fail}$ and $\vec{\alpha}_{sig/bkg,pass/fail}$.

The signal pdf is obtained starting from a simulated Drell-Yan sample template where the GSF electron candidates are matched to the generated electrons ($\Delta R < 0.3$). This distribution is then convoluted with a gaussian function with a mean μ and a width σ (“gmean” and “gsigma” in figures 5.14 and 5.15) to take into account a possible smearing of the data compared to the simulation.

The choice of using a template rather than a theoretical distribution (such as a Breit-Wigner distribution convoluted with a Crystal Ball) has several motivations. First, the kinematic cuts on both the tag and the probe (E_T and η) can induce distortions of the mass spectrum. The low mass tail is for instance always suppressed by the E_T cuts. The Drell-Yan high mass tail is enhanced for $E_{T,probe} \approx 50$ -100 GeV, as illustrated in the middle plots of figure 5.11. Describing this effect with an analytical function is not straightforward. The second reason is that it avoids to give too much freedom to the signal function and limits the risk for the fit to incorporate background events in the signal region. The templates and the parameters of the gaussians used in the functions $f_{sig,pass}$ and $f_{sig,fail}$ are not constrained to be equal: as can be seen from figure 5.11, the peak of the distribution for failing probes is broader than the passing probe distribution.

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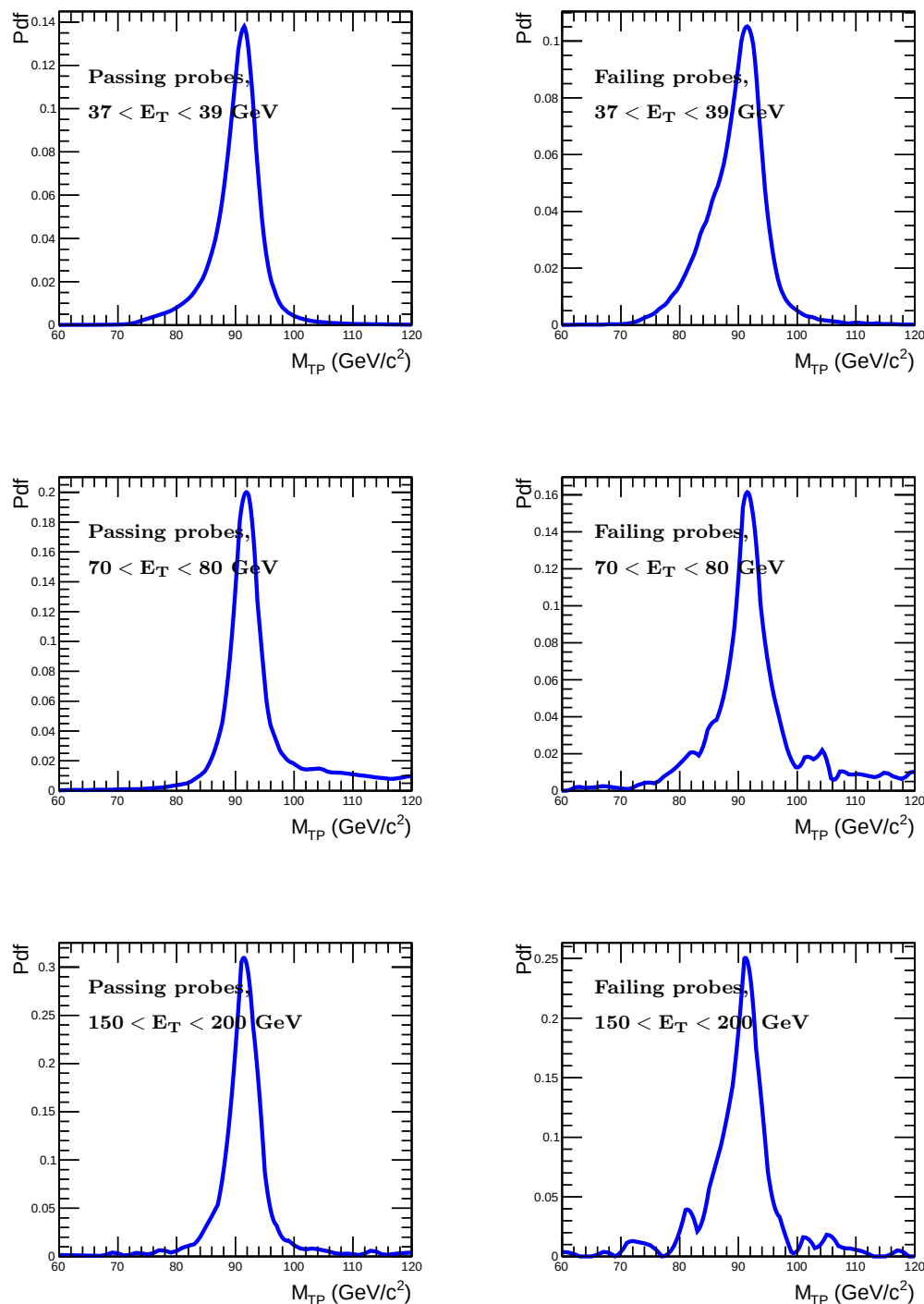


Figure 5.11: The template distributions of the tag and probe invariant mass for passing (left) and failing (right) probes used to build the signal pdf, for three different E_T ranges for the probe : $37 < E_T < 39$ GeV (up), $70 < E_T < 80$ GeV (middle) and $150 < E_T < 200$ GeV (down) for pairs where the probe is reconstructed in the barrel.

	Barrel	Endcaps
$35 \text{ GeV} < E_{T,probe} < 50 \text{ GeV}$	60	60
$50 \text{ GeV} \leq E_{T,probe} < 75 \text{ GeV}$	40	40
$75 \text{ GeV} \leq E_{T,probe} < 100 \text{ GeV}$	40	30
$100 \text{ GeV} \leq E_{T,probe} < 150 \text{ GeV}$	20	20
$150 \text{ GeV} \leq E_{T,probe} < 200 \text{ GeV}$	20	10
$200 \text{ GeV} \leq E_{T,probe}$	20	10

Table 5.4: Number of bins used in the fit of the invariant mass distribution for different probe η and E_T regions.

The background pdf is described by an exponential function multiplied by an error function¹:

$$f_{bkg}(M) = e^{-\alpha \cdot M} \cdot \left(1 + \text{Erf} \left(\frac{M - p_1}{p_2} \right) \right) \quad (5.8)$$

This second factor introduces a threshold effect to model the suppression of low mass events due to the E_T cut illustrated in figure 5.12. The parameters α, p_1, p_2 (denoted “bkgSlope”, “bkgpar1” and “bkgpar2” respectively in figures 5.14 and 5.15) in the functions $f_{bkg,pass}$ and $f_{bkg,pass}$ are not constrained to be equal.

Depending of the probe E_T and η region, the available statistics in both data and simulation strongly vary. The number of bins used for the fit to the invariant mass distribution is therefore adapted to the available statistics and is shown in table 5.4. Figure 5.13 shows the high E_T electron selection efficiency as a function of E_T for barrel and endcaps electrons separately for a simulated Drell-Yan sample and for the data. The same fitting procedure was applied to both data and simulation. The scale factor is shown on the bottom part of the plot and is observed to be flat and close to 1. The errors drawn for the efficiencies are those obtained directly from the fit. No systematic uncertainty (on the background shape,...) is applied. Some examples of the fits on the data are shown in figure 5.14 and 5.15. The efficiencies presented in figure 5.13 are in good agreement within 1% with the ones presented in section 5.2.1 (figure 5.3).

¹ The error function, Erf, is defined as $\text{Erf}(x) = \frac{2}{\sqrt{\pi}} \int_0^x e^{-t^2} dt$. It is a smooth step function going from -1 ($x = -\infty$) to +1 ($x = \infty$).

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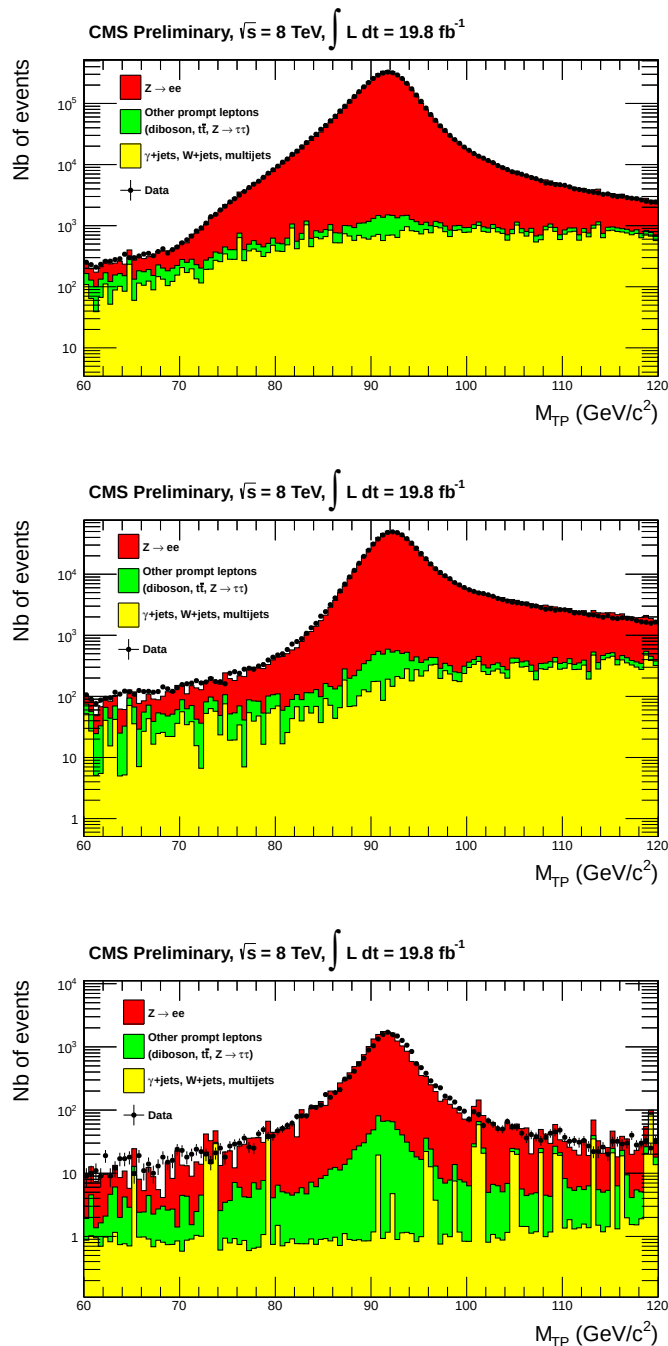


Figure 5.12: The invariant mass of the tag and probe pair for different $E_{T,probe}$ cuts (from top to bottom: $E_T > 35$ GeV, $E_T > 50$ GeV, $E_T > 100$ GeV), for tag and probe electrons emitted in the barrel. The cut on $E_{T,tag}$ is 35 GeV in all the cases. The asymmetric tails induced by the cuts are particularly visible for the multijets background.

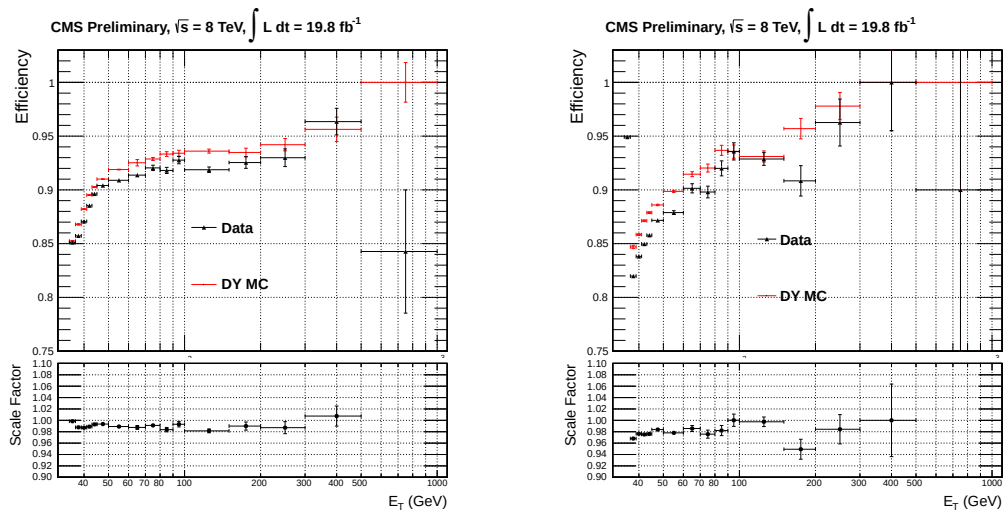


Figure 5.13: High E_T electron selection efficiency and scale factor obtained as a function of E_T using a fitting method to subtract the background. The errors drawn are obtained from the fit.

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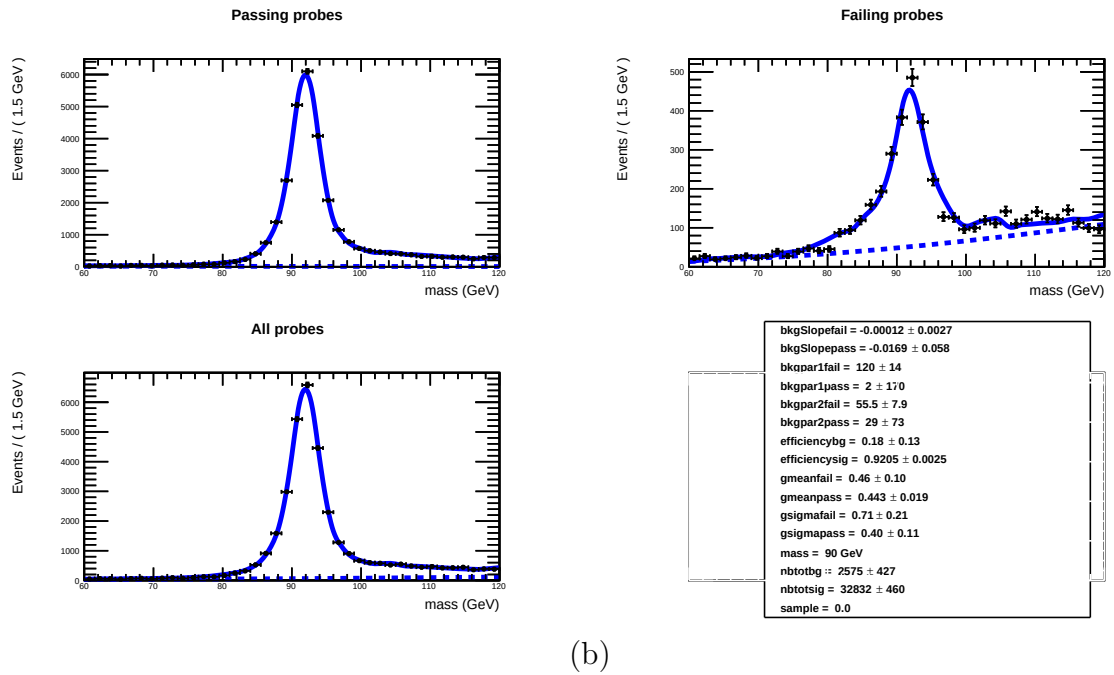
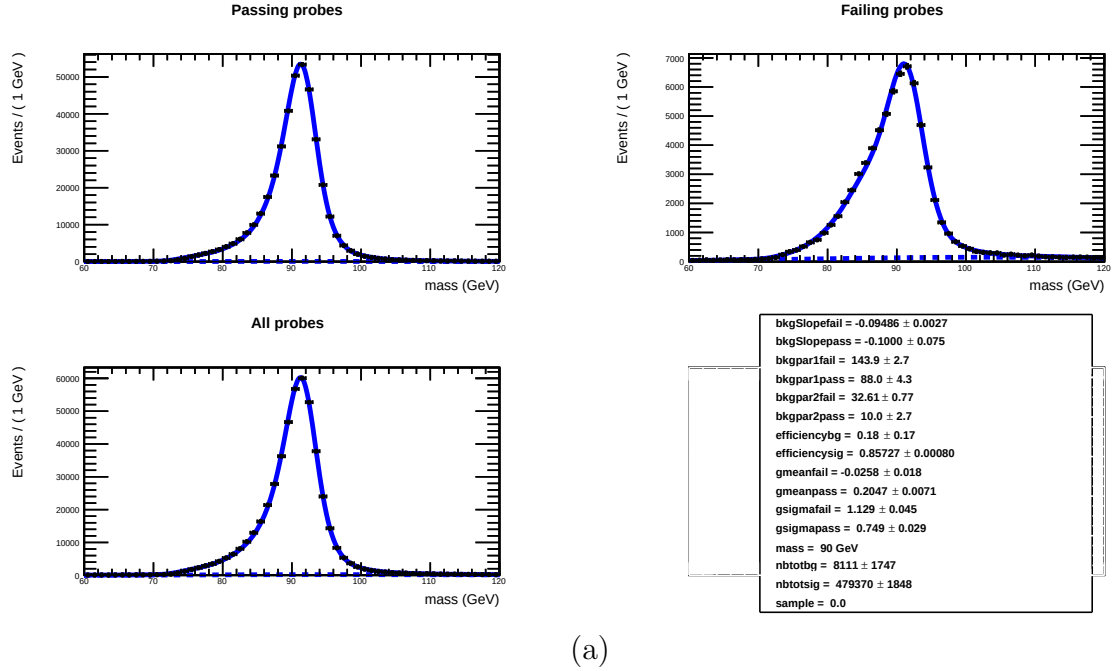


Figure 5.14: Example of fits on the data for probes reconstructed in the barrel in two different E_T ranges : $37 < E_T < 39$ (a) and $70 < E_T < 80$ (b).

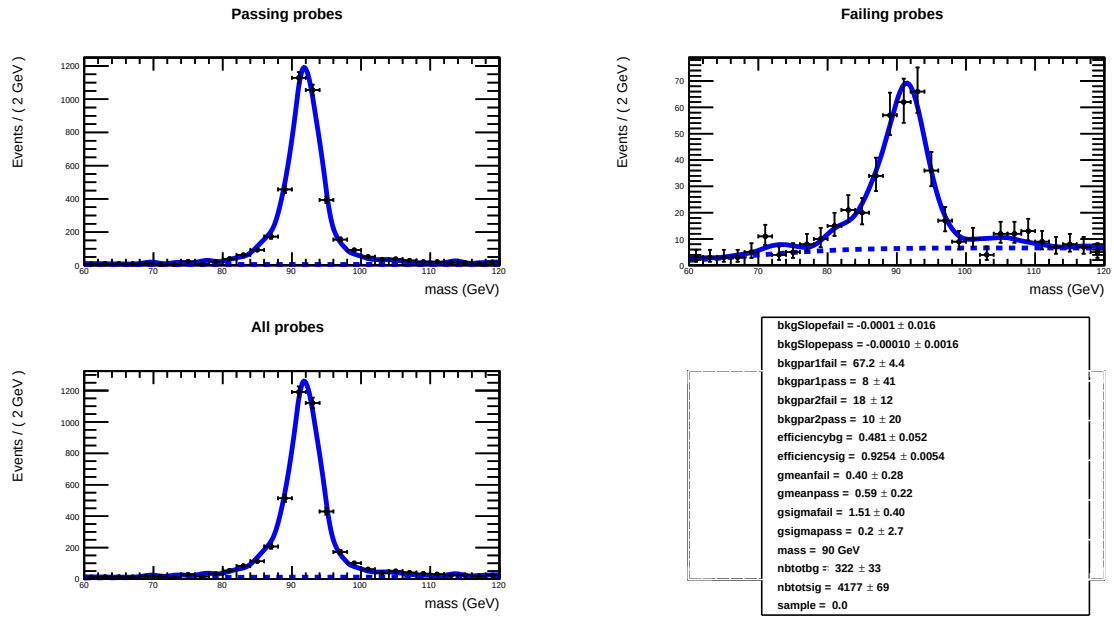


Figure 5.15: Example of fits on the data for probes reconstructed in the barrel and in the range $150 < E_T < 200$.

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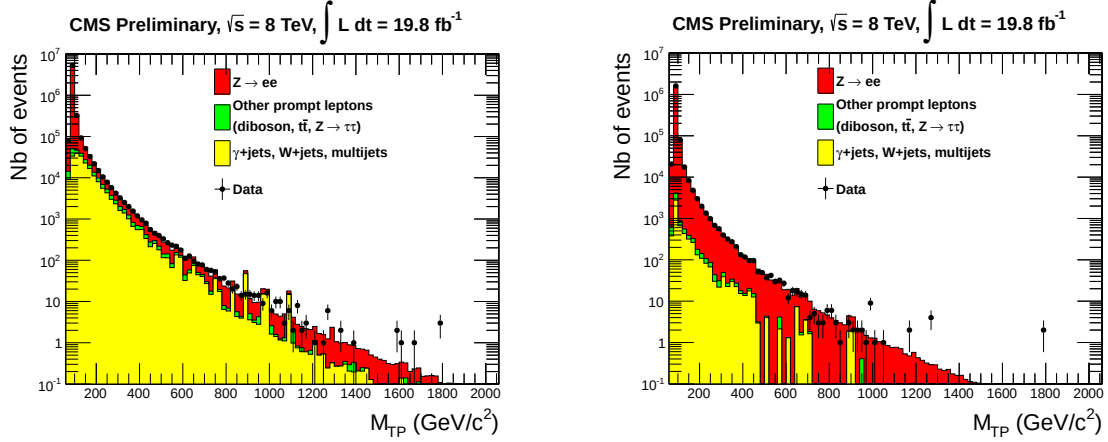


Figure 5.16: Invariant mass spectrum for tag and probe pairs with the selection used in section 5.2.1 (left) and 5.3 (right). Events with two tags are counted twice.

5.3 Results at high mass

The Z peak region allows one to work with a very pure sample of events and thanks to the high statistics, it is even possible to probe electrons with E_T of several hundreds of GeV. However, for high E_T electrons, the topology of the event is very specific with at least one high E_T jet to compensate the transverse impulsion of the on-shell Z boson. It is therefore useful to check the measurement above the Z mass region.

The tag and probe invariant mass distribution is shown in left part of figure 5.16 where, as in section 5.2.1, the tag passes the high E_T electron selection and is emitted in the barrel and the probe is a GSF electron candidate. The background is very important and represents more than 50% of the pairs, as indicated in the first column of table 5.6 which gives the relative yields of the various processes for pairs with $M_{TP} > 200$ GeV. In order to reduce this background, several cuts on the tag and on the event are applied. The list of the extra cuts on the events is shown in table 5.5. The contribution from events where the tag is a jet faking an electron can be sizeably reduced by tightening the conditions on the tag. In order to remove W+jets, Z+jets and γ +jets where the tag is an electron or a converted photon and the probe comes from the jets, cuts on the events are also applied. The two electrons are required to be well balanced in E_T : the difference of E_T over their sum has to be smaller than 0.1 (that is their E_T should match within $\approx 20\%$). This rejects events where particles other than the two electrons carry a high transverse momentum. The missing transverse energy is also required to be smaller than 20 GeV, which allows to significantly decrease the contribution from processes such as W+jets or $t\bar{t}$ where one or several neutrinos are emitted.

The distributions of these two variables can be seen in figure 5.17 for tag and probe pairs with $M_{TP} > 200$ GeV.

As in section 5.2.1, the contribution of the background from prompt leptons is estimated from simulation. However, a different technique than in previous sections is used to

Tag selection	
Variable	Cut value
High E_T electron selection	true
$ \eta_{SC} $	< 1.442
$ \Delta\phi_{in} $	< 0.01
$ \Delta\eta_{in} $	< 0.002
H/E	< 0.02
Track Isolation	$< 1 \text{ GeV}/c$
E/p	< 2
Cuts on the event	
Missing transverse energy	$< 20 \text{ GeV}$
$\left \frac{E_{T,tag} - E_{T,probe}}{E_{T,tag} + E_{T,probe}} \right $	< 0.1
# GSF with $E_T > 10 \text{ GeV}$	$= 2$

Table 5.5: Selection criteria on the tag and on the event for the tag and probe study at high mass.

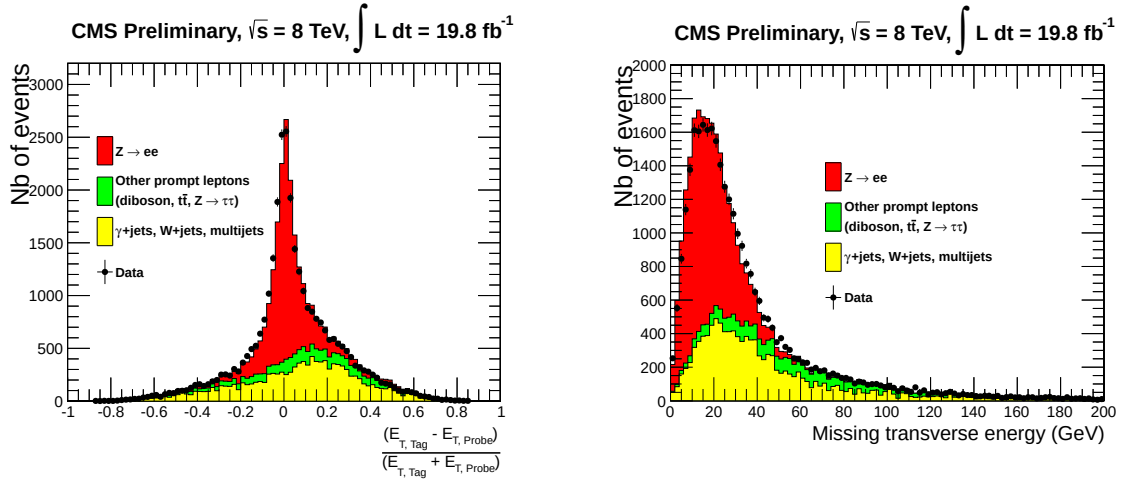


Figure 5.17: Distribution of $\frac{E_{T,tag} - E_{T,probe}}{E_{T,tag} + E_{T,probe}}$ (left) and the missing transverse energy (right) for tag and probe pairs with $M > 200 \text{ GeV}$.

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estimate the contribution from events where the probe is a jet. Because of its small value, the fake rate for this tag selection is in practice hard to measure and subject to a large uncertainty: a small contamination by “real” electrons in the sample used to measure this fake rate can have a large impact on the result.

An alternative way to estimate the contribution of multijets events and also of W+jets and γ +jets is to use the fact that the tag and probe charges are not correlated for such events, as opposed to pairs from Drell-Yan events or other processes involving two prompt leptons. Pairs where both objects carry an identical charge form therefore a sample dominated by this process and is only contaminated by Drell-Yan or other prompt lepton processes where the charge of one of the leptons is wrongly measured. As shown previously in figures 5.7 to 5.10, the probability for this to happen is small and well described by the simulation and simulation are used to remove this residual contribution. The total contribution of processes involving at least one jet can therefore be estimated by taking twice the corrected number of pairs where the two objects carry the same charge.

Second column of table 5.6 indicates the yields after all the extra cuts listed in table 5.5. The set of cuts reduces the number of Drell-Yan events by a factor of three but the sample purity is now over 90%.

	Before additional cuts	After additional cuts
Number of events (data)	55510	7595
Drell-Yan	36.4%	91.1%
Other prompt leptons	10%	1.9%
Jets	53.6%	7.0%

Table 5.6: Number of events in data and relative fractions of the different processes obtained from simulation for tag and probe events with $M_{TP} > 200$ GeV before and after the additional cuts on the tag and on the event.

The efficiency as a function of the tag and probe mass after the additional cuts is shown in figure 5.19. Despite the strong cuts, the background contamination is still present, especially when the probe lies in the endcaps. However a good agreement is observed for the whole spectrum. This is quantified in table 5.18 where the efficiency for pairs with $M_{TP} > 200$ GeV is shown for the data and the simulation. The scale factor, defined as in section 5.2.1 is also computed and agrees nicely with the numbers presented in table 5.3.

	Barrel, $E_T \geq 35$ GeV	Endcaps, $E_T \geq 35$ GeV
Efficiency (Drell-Yan)	$0.943 \pm 0.003(stat.)$	$0.925 \pm 0.006(stat.)$
Efficiency (Drell-Yan + BG)	$0.896 \pm 0.004(stat.)$	$0.826 \pm 0.007(stat.)$
Efficiency (Data)	$0.880 \pm 0.004(stat.)$	$0.809 \pm 0.008(stat.)$
Efficiency (Data BG corr.)	$0.927 \pm 0.004(stat.)$	$0.908 \pm 0.008(stat.)$
Scale factor	$0.982 \pm 0.005(stat.) \pm 0.019(syst.)$	$0.980 \pm 0.010(stat.) \pm 0.040(syst.)$

Figure 5.18: High E_T electron selection efficiency for electrons reconstructed in the barrel (left) and endcaps (right), obtained with tag and probe pairs with $M_{TP} > 200$ GeV and passing the selection indicated in table 5.5, for a Drell-Yan sample, on a Drell-Yan and background sample, and on the 2012 data before and after background subtraction. The last row gives the scale factor, taken as the ratio of the efficiency measured in data after background subtraction to the one measured in the Drell-Yan sample. The systematic error on the scale factor is obtained by varying the expected background contribution by 40%.

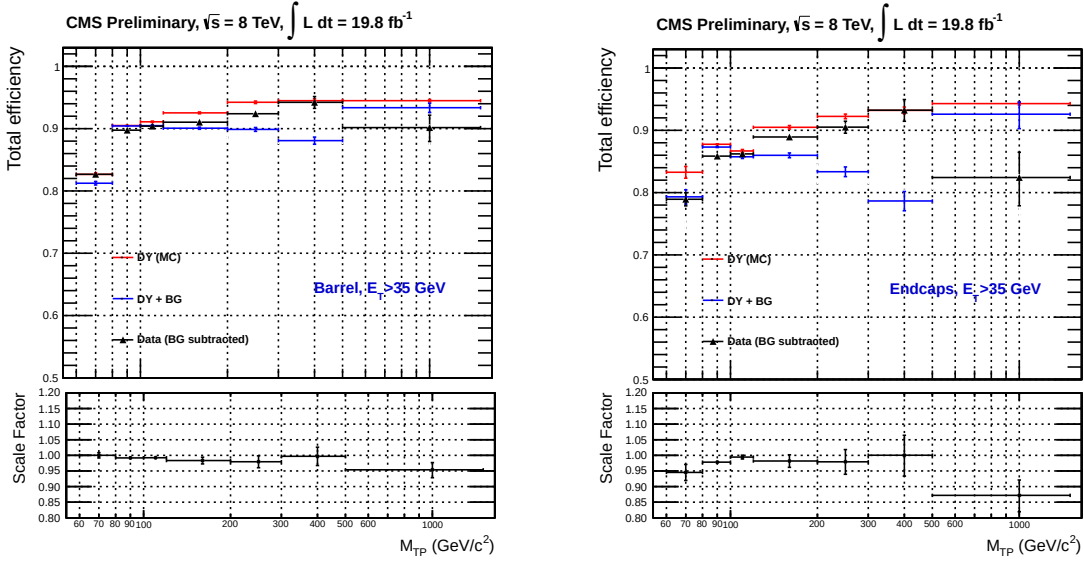


Figure 5.19: High E_T electron selection efficiency and scale factor as a function of the tag and probe invariant mass for probe electrons with $E_T > 35$ GeV with the cuts described in table 5.5. Barrel (left) and Endcaps (right) efficiencies are shown separately. The efficiency is shown for a Drell-Yan only sample (red), for a Drell-Yan + background sample (blue) and for the data before (dots) and after (triangles) background subtraction. Only statistical errors are drawn for the efficiencies. Both statistical and systematic uncertainties are drawn for the scale factor.

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Name	Comments
L1_SingleEG20 or L1_SingleEG22	1 e/γ candidate with $E_T > 20$ or 22 GeV
HLT_Ele27_WP80	1 electron object satisfying: $E_T > 27$ GeV $ \Delta\eta_{in} < 0.007$; $ \Delta\phi_{in} < 0.06$ $H/E < 0.1$ $\sigma_{in\eta} < 0.01$ Tracker isolation $< 0.05 \times E_T$ ECAL isolation $< 0.15 \times E_T$ HCAL isolation $< 0.1 \times E_T$ $1/E - 1/p < 0.05$

Table 5.7: The L1 (first line) and HLT (second line) triggers used for the measurement of the GSF electron candidate reconstruction and trigger efficiencies. Only values for barrel electrons are quoted since the tag selection imposes the electron to be emitted in this region.

5.4 Reconstruction and trigger efficiency

Sections 5.2 and 5.3 focus on the measurement of the efficiency of the high E_T electron selection, which starts with a GSF electron candidate. The measurement is performed with data collected with the trigger described in table 4.2.

The other efficiencies appearing in equation 5.1 needs however also to be measured. This is the topic of this section where results are presented for ϵ_{GSF} and $\epsilon_{trigger}$ (ϵ_{SC} was measured to be 100% [9]).

The measurements presented here use the tag and probe technique and is based on a fit to the invariant mass distribution of tag and probe pairs at the Z pole, as in section 5.2.2. The tag selection is identical but the probe is different for the reconstruction efficiency.

In order to avoid any bias in the measurements, the trigger used here requires only one electron object, that will be matched with the tag. It is detailed in table 5.4. At the L1, the E_T threshold used was dependent on the data taking conditions: for high instantaneous luminosities, a cut at 22 GeV was needed to maintain the L1 trigger rate at an acceptable level. For lower luminosities, the cut is lowered down to 20 GeV.

5.4.1 GSF electron candidate efficiency

Here, the probe is taken to be a supercluster without any matching requirement with a track, and the efficiencies are computed with respect to this preselection. The probability for an electron emitted in the ECAL acceptance to form a supercluster is therefore the only step not measured in the data. It is estimated from simulation and found to be around 99% for electrons with $E_T > 35$ GeV. The efficiency of the GSF electron candidate algorithm is shown as a function of the supercluster E_T in figure 5.20, for probes reconstructed in the barrel and endcaps separately. Except in the E_T region below 30 GeV in the barrel, an excellent agreement is observed between the data and the simulation. The small variation

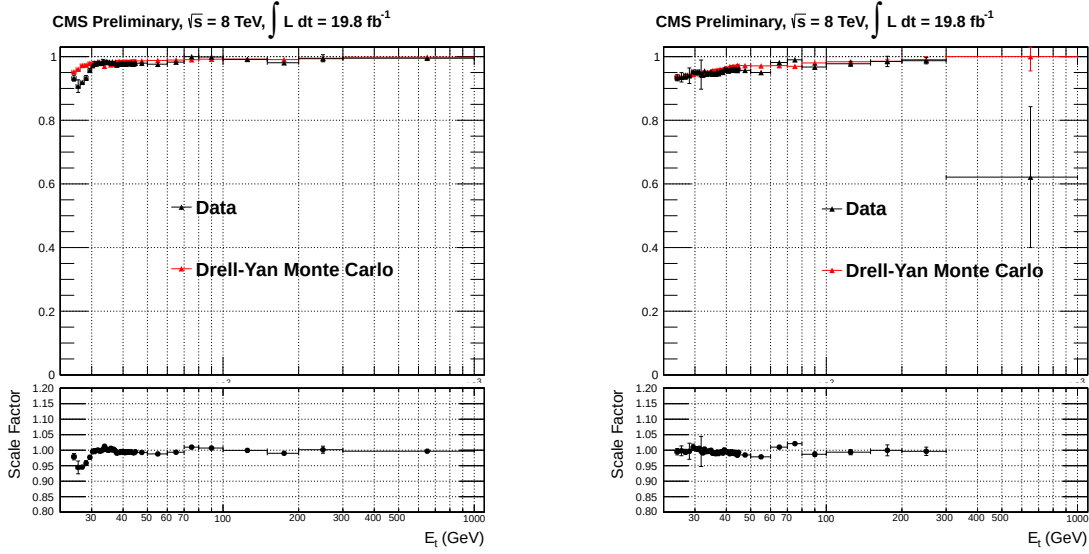


Figure 5.20: Efficiency of the GSF electron candidate algorithm as a function of the supercluster E_T , for probes reconstructed in the barrel (left) and endcaps (right). The errors drawn are obtained from the fit.

of 1% observed between data and simulation in a few bins are thought to come from the uncertainty on the signal and background shapes.

5.4.2 Trigger efficiency

The HLT trigger studied here is made of two identical electron “legs” and based on a single L1 object. The efficiency of the L1 object and of a leg of the HLT trigger are computed with respect to GSF electron candidates and are presented as a function of E_T in figure 5.21 and 5.22 respectively.

For the L1, a strong discrepancy between the data and the simulation is observed below 30 GeV of E_T . This is mainly due to the fact that the energy measurement for the L1 object and the GSF electron candidate are not exactly the same. The crystals used in the L1 tower and in the supercluster are not identical and several energy corrections are applied to the GSF electron candidates. The energy of the L1 object is not very well known and the efficiency of the E_T cut for GSF electron candidate close to the L1 threshold is therefore hard to simulate. For GSF electron candidates with $E_T > 35$ GeV, the effect is much reduced and the efficiency in data is above 95%. The discrepancy vanishes completely for $E_T > 50$ GeV. Since two electrons are finally required at the HLT level and that each of them can fire the L1 trigger, the discrepancy between data and simulation for the L1 efficiency is negligible and the L1 efficiency close to 100% in both.

The HLT efficiency shows a behaviour comparable to the L1. For E_T values above 50 GeV, the data and simulation are in good agreement and tend to 98-99%. At lower values, a turn-on curve is observed due to the 33 GeV E_T cut in the HLT. The drop is present in both data and simulation but its position and its width are not well modeled. Although

5. MEASUREMENT OF THE HIGH E_T ELECTRON SELECTION, RECONSTRUCTION AND TRIGGER EFFICIENCIES IN THE 2012 DATA

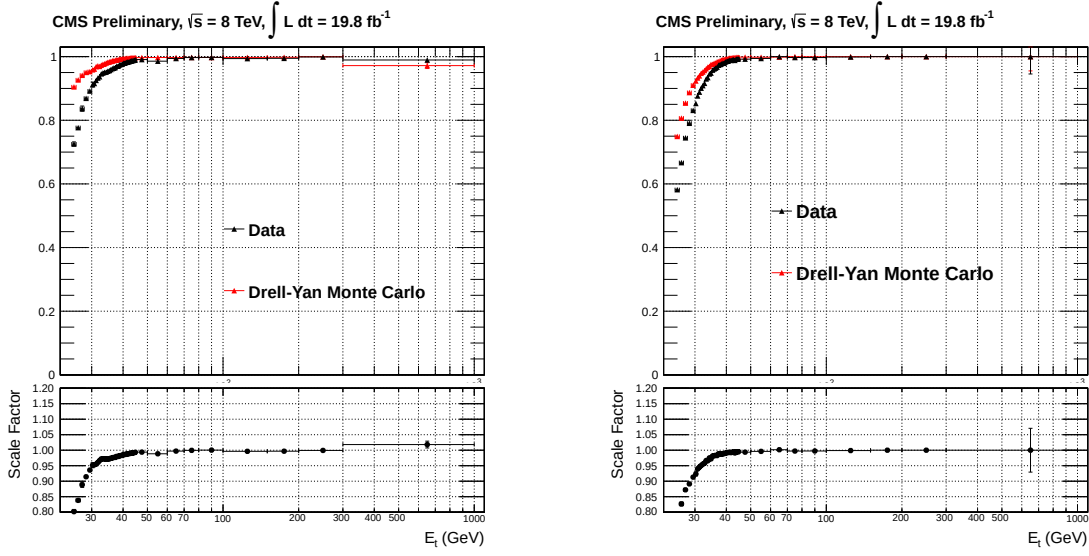


Figure 5.21: Efficiency of the L1 trigger as a function of the GSF electron candidate E_T , for probes reconstructed in the barrel (left) and endcaps (right) The errors drawn are obtained from the fit.

this feature has no effect for electron in the signal region of the present analysis, it can significantly alter measurements using electrons of $E_T \approx 35$ GeV. As will be detailed in chapter 6, limits will be set on the ratio of the cross section production of a new heavy dielectron resonance over the measured Z resonance cross section. A good description of the trigger efficiency at $E_T \approx 35$ GeV is therefore needed. The final selection used for the search for new heavy dielectron resonances requires two GSF electron candidates passing the high E_T electron selection. In order to accurately describe the trigger, simulated events are not required to pass the trigger such as described in the simulation but are rather reweighted so that the two electrons follow the trigger efficiency curve measured in the data.

The efficiency can be factorized in two parts: the E_T threshold turn-on and the asymptotic trigger efficiency for electrons with E_T in the plateau of the turn-on curve.

To measure the turn-on, a parametric function is fit to the data using a sample of tag and probe pairs where the tag and the probe are GSF electron candidates passing the high E_T electron selection and where one of them is reconstructed in the barrel and the other in the endcaps. Since one will apply the reweighting only to electrons passing the high E_T electron selection, the turn-on curve must indeed be measured with respect to probes passing this selection. As a consequence, this efficiency differs from $\epsilon_{trigger}$ in equation 5.1 or in figure 5.22. This selection also ensures that the contamination of the sample by fake electrons is negligible. No background subtraction is therefore performed here and the efficiency is simply obtained by counting the number of passing pairs over the total number of pairs. The efficiency is shown in figure 5.23 as a function of the probe E_T .

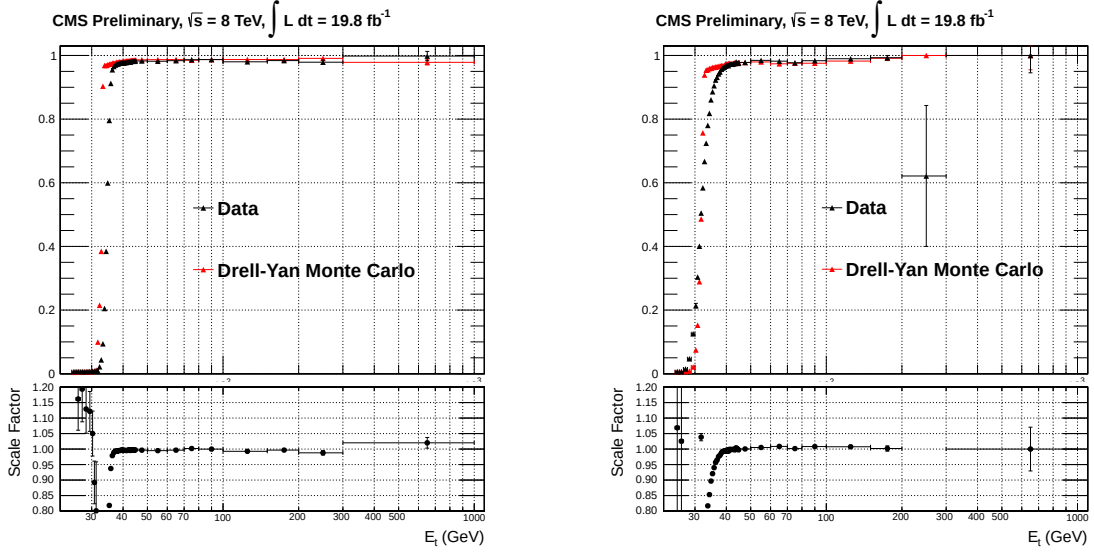


Figure 5.22: Efficiency of a leg of the HLT trigger as a function of the GSF electron candidate E_T , for probes reconstructed in the barrel (left) and endcaps (right). The errors drawn are obtained from the fit.

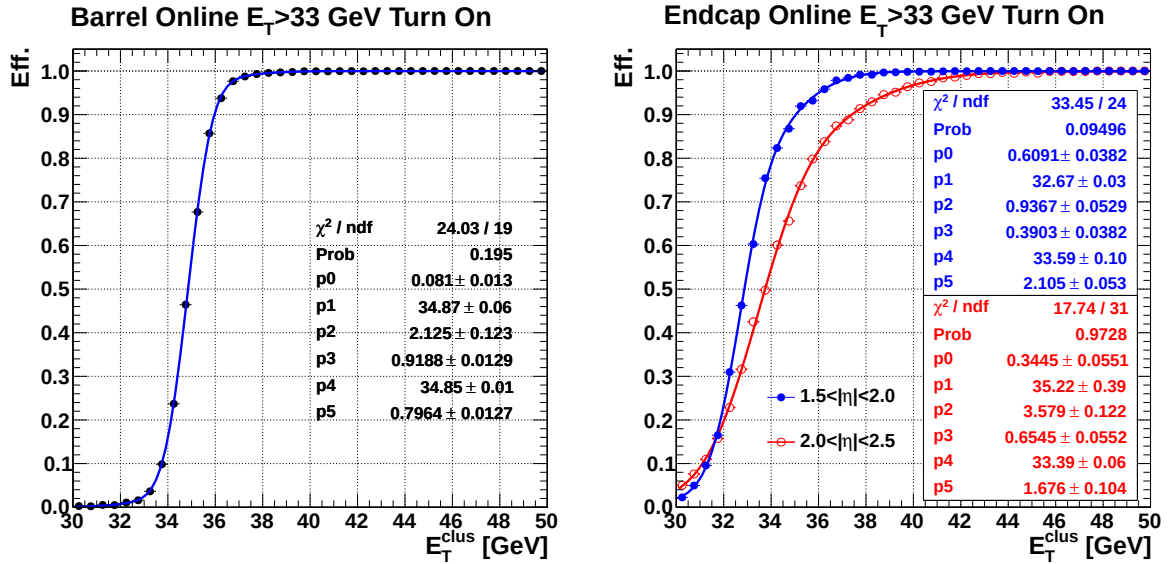


Figure 5.23: Efficiency of the HLT electron leg for GSF electron candidates passing the high E_T electron selection in the data using the tag and probe method. The efficiency is shown separately for electrons emitted in the barrel (left) or in the endcaps (right). For the endcaps, electrons are subdivided into two regions: $|\eta| < 2.0$ (blue) and $|\eta| > 2.0$ (red). The results of a fit to the parametric function given in equation 5.9 is also shown.

5. MEASUREMENT OF THE HIGH E_T ELECTRON SELECTION, RECONSTRUCTION AND TRIGGER EFFICIENCIES IN THE 2012 DATA

A parametric function consisting of the sum of two error functions:

$$f(E_T) = 0.5p_0 \left(1 + \text{Erf} \left(\frac{E_T - p_1}{\sqrt{2}p_2} \right) \right) + 0.5p_3 \left(1 + \text{Erf} \left(\frac{E_T - p_4}{\sqrt{2}p_5} \right) \right) \quad (5.9)$$

is fit to the data and the estimation of the fit parameters is shown in figure 5.23. Different fits are performed for probes in the barrel and in the endcaps. In the latter case, probes are furthermore separated into two categories: $|\eta| < 2.0$ (blue) and $|\eta| > 2.0$ (red).

The weight ω applied to a simulated event with two GSF electron candidates passing the high E_T electron selection to emulate the trigger turn-on is therefore:

$$\omega = f(E_{T,1}, \eta_1) \times f(E_{T,2}, \eta_2) \quad (5.10)$$

The asymptotic trigger efficiency for electrons passing the high E_T electrons selection is measured in the same way once the plateau is reached and found to be $99.65 \pm 0.02\%$ and $99.61 \pm 0.02\%$ for GSF electron candidates emitted in the barrel and in the endcaps respectively.

5.5 Summary

In this section, the efficiencies of the various steps of the final electron selection used in the search for new heavy resonances were studied. The efficiencies were measured both in data and simulation using the tag and probe method. A special attention was paid to the behaviour of the high E_T electron selection efficiency at very high E_T (up to ≈ 300 - 500 GeV). The ratio of the data over simulation measurements lies between 0.98 and 1.00 in the whole E_T range considered. Several complementary techniques were used and lead to measurements in agreement with each other. Given these results, a systematic uncertainty of 2% is assigned to the evolution of the efficiency scale factor to the signal region (E_T up to 1 TeV) for electrons emitted in the barrel. The reduced statistics in the endcaps at high E_T and small discrepancies observed in that region between data and simulation leads to the assignment of a more conservative uncertainty of 4%.

The GSF electron candidate reconstruction and the L1 and HLT triggers were found to be extremely efficient and the data over simulation disagreement to be below the systematic uncertainties indicated for the high E_T electron selection. No corrections are therefore applied to the simulation predictions. An exception is the turn-on curve due to the E_T threshold of the HLT trigger which is poorly described. A parametrized function was fit to the data and is applied to the simulation to emulate the trigger turn-on.

Chapter 6

Search for new heavy dielectron resonances with the data collected by CMS at $\sqrt{s} = 8$ TeV

The search for new heavy resonances decaying into a dilepton pair at hadron colliders is a topic of interest since many years. Before the LHC era, the strongest limits on the production cross sections of such phenomena were obtained by the CDF and D0 collaborations at the Tevatron [43, 44, 45]. This search benefits from the high center of mass energies that can be reached at the LHC together with the clean final state of the signal. It therefore provides a clear way to observe new physics at the TeV scale.

This chapter presents the results of the search for new heavy resonances decaying into a dielectron pair using the proton-proton collisions at $\sqrt{s} = 8$ TeV recorded by the CMS detector in 2012. They have been the subject of a Physics Analysis Summary [3] released by the CMS collaboration and a paper which interprets the observed dielectron and dimuon invariant mass spectrum in terms of searches for both resonant or non resonant signals is currently in preparation [4]. This paper follows previous ones based on the analysis of the data recorded at $\sqrt{s} = 7$ TeV in 2010 [5] and on the combination of the analysis of the data recorded at $\sqrt{s} = 7$ TeV in 2011 and of the first part of the 2012 data [7]. A similar search was also performed by the ATLAS Collaboration [120, 121] which led to comparable results to those presented here.

Section 6.1 gives the datasets used in this analysis and presents the event selection and the product of the signal acceptance times efficiency. In section 6.2, dielectron pairs at the Z pole are used to compare the performance of the electron energy measurement in the data and simulations. Simulations are then used to estimate the dielectron mass resolution at high mass. The various backgrounds leading to the reconstruction of a high mass dielectron pair are detailed in section 6.3 and estimated or cross checked with data techniques when possible. The final dielectron mass spectrum measured in the data is shown in section 6.4 and compared to the Standard Model predictions. The statistical interpretation is presented in section 6.5. Since no significant deviation to the Standard Model prediction is observed, upper exclusion limits are set on the production cross section of new heavy narrow resonances. Those limits are converted in a lower limits on the mass of new particles predicted by various models of new physics.

6.1 Datasets and event selection

The data analyzed in this thesis correspond to the complete dataset recorded in 2012 by the CMS detector. They are listed in table 2.3. The total integrated luminosity is $19.7 \pm 0.5 \text{ fb}^{-1}$ [61]. Various triggers were used for the complete analysis, the main one being the trigger described in table 4.2 which is used for the final dielectron mass spectrum measurement.

6.1.1 Event selection

The event selection requires two GSF electron candidates (*cf.* section 4.1.3) with a transverse energy above 35 GeV passing the high E_T electron selection described in table 4.4 and whose supercluster is reconstructed within the tracker detector acceptance ($|\eta_{SC}| < 2.5$) excluding those in the ECAL barrel and endcaps transition region ($1.442 < |\eta_{SC}| < 1.56$). If more than one pair of electrons satisfy these criteria, only, the pair formed by the two electrons with the highest E_T is considered. The event must also pass the trigger described in table 4.2.

At least one of the electron must in addition be reconstructed in the ECAL barrel. This condition is motivated by the fact that the background contribution from events where two jets are reconstructed as fake electrons in the opposite endcaps (*cf.* section 6.3.3) is important whereas this configuration is kinematically disfavoured for the signal at high mass. Because of the different performance of the electron reconstruction and selection in the ECAL barrel and endcaps, events are separated between those where both electrons are reconstructed in the barrel (barrel-barrel category) and those where one of them is reconstructed in one of the endcaps (barrel-endcaps category).

No requirement on the sign of the two electrons is made as the probability for an electron to be assigned the wrong charge is sizeable at high E_T and such a condition would lead to a 5% loss of efficiency (*cf.* section 5.2.1.1) for only a reduction of a factor two of the backgrounds where the two reconstructed electrons are not correlated. As will be shown in section 6.3, those backgrounds consist of events where at least one of the electron is a fake and represent only a couple of percents of the total background yield. The signal loss of the opposite-sign condition is therefore not justified.

No additional conditions (on the number on jets, on the missing transverse energy, on the azimuthal angle between the two electrons,...) are imposed. The reason is that the present search aims to be as model independent as possible and does not make any assumption on the signal except that it produces a narrow dielectron resonance. For models where the new resonance is not produced in association with other particles, the event by event signature of the signal is identical to the one of the Drell-Yan process which represents the leading Standard Model background. Adding extra conditions on the event is therefore not expected to improve the results in that case but would on the other hand make the analysis blind to more exotic scenarios where other particles than the two electrons are present in the final state.

6.1.2 Signal acceptance times efficiency

The probability for a proton-proton collision where two high energy electrons are produced in the final state to be recorded and to pass the full event selection can be split into two parts:

- a kinematical “acceptance” defined as the probability for the two electrons to be reconstructed inside the instrumented area of the detector and to pass the E_T threshold.
- the product of the efficiencies of the reconstruction and selection of the two electrons.

Figure 6.1 shows the cumulated acceptance times efficiency of the various steps of the event selection as a function of the mass of dielectron pairs coming from Drell-Yan simulated events. For masses above 2 TeV, around 90% of them are reconstructed in the detector acceptance. Less than 5% of the dielectrons pairs have both electrons emitted in the ECAL endcaps. The reconstruction of two GSF electron candidates and the trigger generates a loss of about 8% of the events. This value is slightly higher than one would expect by squaring the efficiencies measured in sections 5.4.2 and 5.4.1 ($\approx 5\%$). The difference originates from the existence of a η condition ($|\eta| < 1.442$ or $1.56 < |\eta| < 2.5$) on the GSF electron candidate in figure 6.1 and from the fact that very forward electrons for which the efficiency is smaller than in the rest of the detector are kinematically disfavoured due to the event selection in the measurement presented in sections 5.4.2 and 5.4.1. The requirement that both candidates pass the high E_T electron selection rejects 12% of the events. The total selection acceptance times efficiency for Drell-Yan events is flat for invariant masses above 1 TeV and around 65%.

Since a new signal leading to a heavy dielectron resonance is kinematically similar to high mass Drell-Yan pairs, the selection acceptance times efficiency is expected to be similar. However, it can slightly differ because of two effects. First, the angular distributions of the dielectron pair vary, depending on the spin of the resonant object and on its production mode. Their shapes for spin 0, 1 or 2 resonances are discussed in the next chapter and their theoretical angular differential cross sections are given in table 7.3. This effect can typically lead to variation of the acceptance of a couple of percents for masses above 1 TeV. Secondly, different couplings to up and down quarks give rise to a different contribution from the valence quarks of the proton and will impact the longitudinal boost of the pair which can modify the geometrical acceptance. The possible coupling to gluons for a spin 0 or 2 resonance also affects the longitudinal momentum distribution of the pair. This effect was however checked to be negligible and, for a given spin and production mode, the acceptance can in the cases be assumed to be universal.

Figure 6.2 shows the acceptance times efficiency for the barrel-barrel and barrel-endcaps categories separately for pairs coming from the decay of a spin 0 (top), a spin 1 (middle) or a spin 2 (bottom) object. The spin 1 and spin 2 curves were obtained using Drell-Yan simulated events and a series of samples simulating Randall-Sundrum gravitons with various masses. The spin 0 is obtained using Drell-Yan pairs reweighted such that their angular distribution corresponds to an isotropic decay. Sizeable differences can be observed, especially for masses below 500 GeV. The fractions of barrel-barrel and barrel-endcaps events strongly vary.

6. SEARCH FOR NEW HEAVY DIELECTRON RESONANCES WITH THE DATA COLLECTED BY CMS AT $\sqrt{s} = 8$ TEV

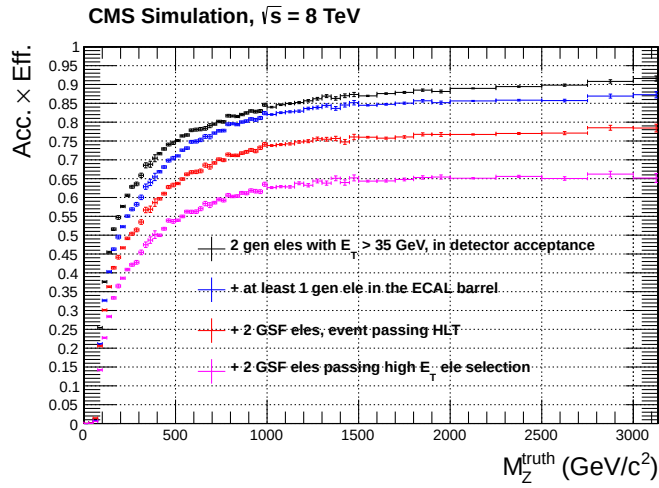


Figure 6.1: Acceptance times efficiency for dielectron pairs coming from simulated Drell-Yan events as a function of the invariant mass of the generated dielectron pair for the various steps of the event selection: The black curve shows the detector acceptance at the generated level and the efficiency of the E_T threshold. The additional condition that at least one of the electrons is emitted in the ECAL barrel leads to the blue curve. The red (magenta) curve requires in addition that the event passes the HLT trigger and the reconstruction of two GSF electron candidates (two GSF electron candidates passing the high E_T electron selection).

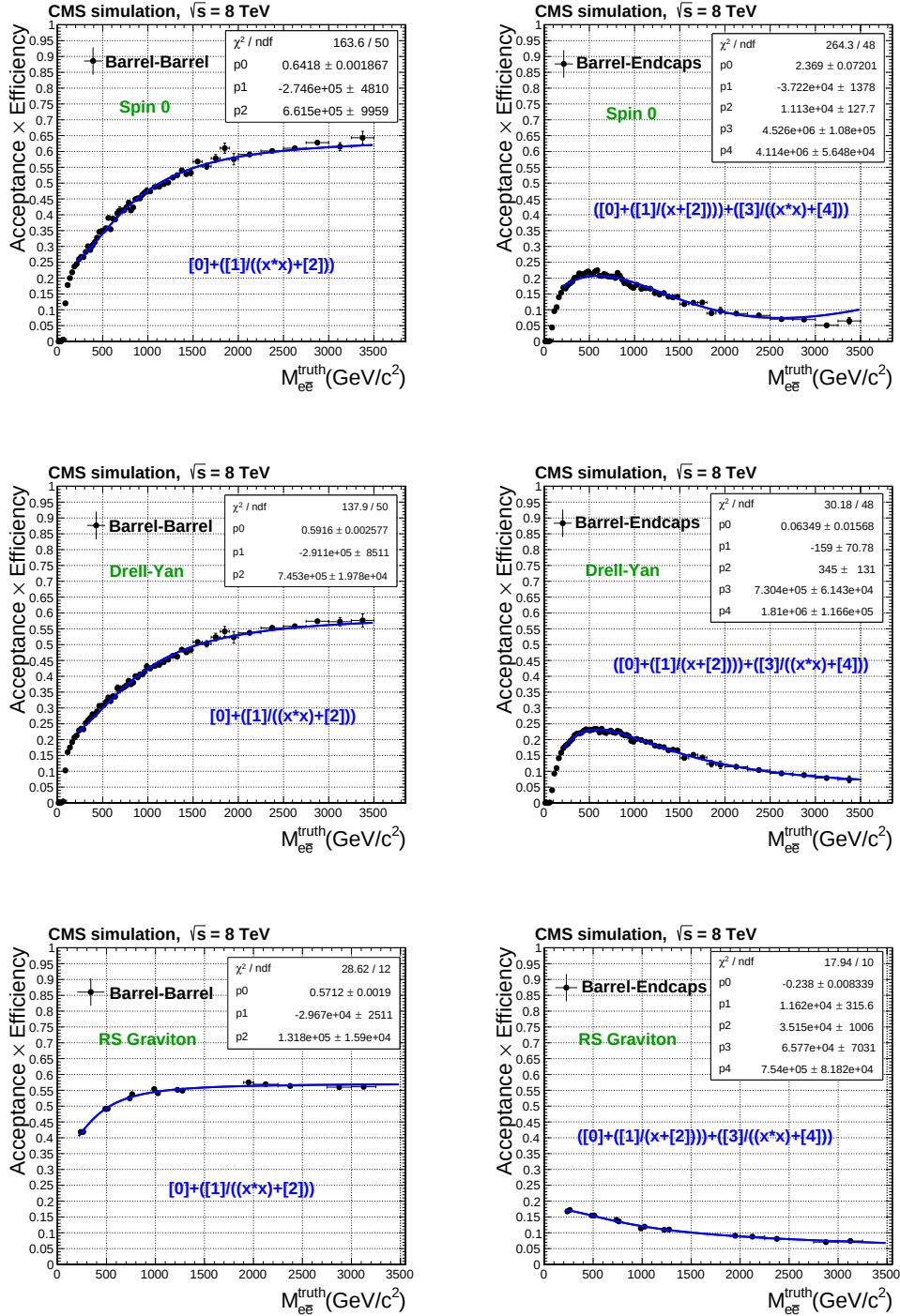


Figure 6.2: Acceptance times efficiency for dielectron pairs coming from the decay of a spin 0 (top), spin 1 (middle) or spin 2 (bottom) object as a function of the dielectron invariant mass. The events are split into two categories: events where both electrons are reconstructed in the ECAL barrel (left), one electron is reconstructed in the ECAL barrel and one in the endcaps (right). The result of a fit to the function described in equations 6.1 is also shown.

6. SEARCH FOR NEW HEAVY DIELECTRON RESONANCES WITH THE DATA COLLECTED BY CMS AT $\sqrt{S} = 8$ TEV

In section 6.5, limits will be set on the production cross section of spin 0, 1 or 2 resonances decaying into a dielectron pair. An input to the limit calculation is the parametrization of the acceptance times efficiency as a function of the mass. A fit to the points in figure 6.2 is therefore performed. A different parametric function was used for each category:

$$p_0 + \frac{p_1}{M^2 - p_2} \quad (\text{barrel} - \text{barrel}) \quad (6.1)$$

$$p_0 + \frac{p_1}{M - p_2} + \frac{p_3}{M^2 - p_4} \quad (\text{barrel} - \text{endcaps}) \quad (6.2)$$

The fit parameters and the resulting curve are shown in figure 6.2.

6.2 Electron energy scale and dielectron mass resolution

In the search for dielectron resonance, a precise measurement of the electron energy is crucial. For narrow resonances, the mass resolution determines the typical range of mass where the signal has a sizeable contribution. A bad estimation of this width can therefore strongly impact the results of the analysis.

The electron energy scale, defined as the average relative shift between the electron energy measured in the simulation and in the data, and the dielectron mass resolution were measured using dielectrons pairs at the Z pole. The expected theoretical shape of a resonance follows a Breit-Wigner distribution. This distribution must be convoluted with the experimental measurement uncertainty, which is often modeled by a Gaussian function. In the case of electrons however, it can happen that a large fraction of the energy is missed because of the radiation of a high energy photon arising from the interaction of the electron with the tracker detector. The supercluster algorithm 4.1.1 is designed to add the contribution of these photons to the electron energy but sometimes fails to recover them.

To model that, a Crystal Ball function is usually used rather than a Gaussian function. It is defined as a truncated Gaussian distribution where the low energy tail is replaced by a decreasing exponential.

In this study, a double sided Crystall Ball function (where both tails are decreasing exponentially) was found to better fit the data and was therefore used. The physical reason for adding a long high energy tail is that in a few cases the supercluster includes energy deposits not associated to the electron. This high energy tail is however smaller than the low energy one.

The dielectron invariant mass distribution obtained from pairs of GSF electron candidates passing the high E_T electron selection was therefore fitted to the convolution of a Breit-Wigner resonance with a double sided Crystall Ball. In order to avoid any distortion of the low mass part of the resonance, the E_T threshold was lowered down to 25 GeV and a dedicated trigger was used. The results are shown separately for the barrel-barrel and barrel-endcaps categories in figure 6.3 for the data and a Drell-Yan simulated sample. The non Drell-Yan contribution to the mass spectrum in the data was found to be negligible

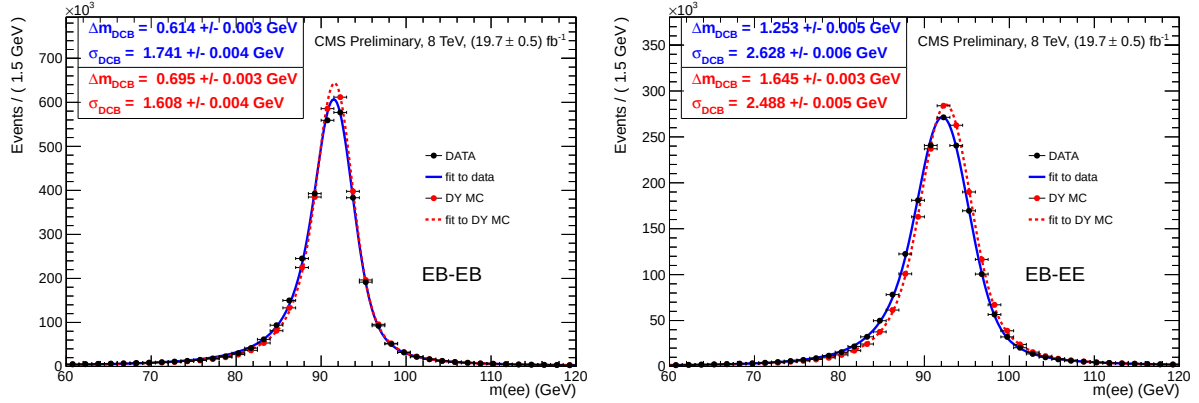


Figure 6.3: Invariant mass distribution of dielectron pairs in the mass range 60 to 120 GeV for the barrel-barrel (left) and barrel-endcaps (right) categories in the data (black) and in Drell-Yan simulated events (red). A fit to a Breit Wigner function convoluted with a double sided Crystall Ball function is performed for both the data (blue) and the simulations (red). The estimated width of the Crystall Ball σ_{CB} and the difference between the peak of the function and the Z mass (91.2 GeV) are also shown. In order not to distort the distribution, the electron E_T cut used here is 25 GeV instead of 35 GeV.

region	$\sigma_{DCB}(data)$	$\sigma_{DCB}(simulation)$	$\sigma_{DCB}(extra)$
barrel-barrel	1.96 ± 0.00	1.74 ± 0.02	0.92 ± 0.02
barrel-endcaps	3.07 ± 0.01	2.71 ± 0.02	1.43 ± 0.02

Table 6.1: The dielectron mass resolution for electron pairs in the Z peak region measured in the data and in the simulations for pairs where both electrons are reconstructed in the ECAL barrel and where one of them is reconstructed in one of the endcaps. The last row shows the quadratic difference between the data and simulation values. All the values are expressed in GeV.

and no background subtraction was therefore performed. A sizeable difference between the data and the simulation is seen: the distribution is broader and slightly shifted to low mass values in the data. An electron energy scale, defined as the difference between the position of the peak in the data and in the simulations divided by the theoretical Z mass value (91.2 GeV) is measured to be $-0.09 \pm 0.00 \%$ for the barrel-barrel category and $-0.43 \pm 0.01 \%$, which is smaller than the mass resolution given by the width of the double sided Gaussian function and whose values in the data and in the simulations are shown in table 6.1. The last row of the table shows the value of $\sigma_{DCB}(extra)$ defined as the quadratic difference between the data and simulation values.

As mentioned in section 2.2.5.1, the ECAL energy resolution is the quadratic sum of three terms which evolve differently with the energy (equation 2.6).

Figure 6.4 studies the evolution of the dielectron mass resolution as a function of the mass of the pair. As expected, its behaviour is similar to the electron energy resolution

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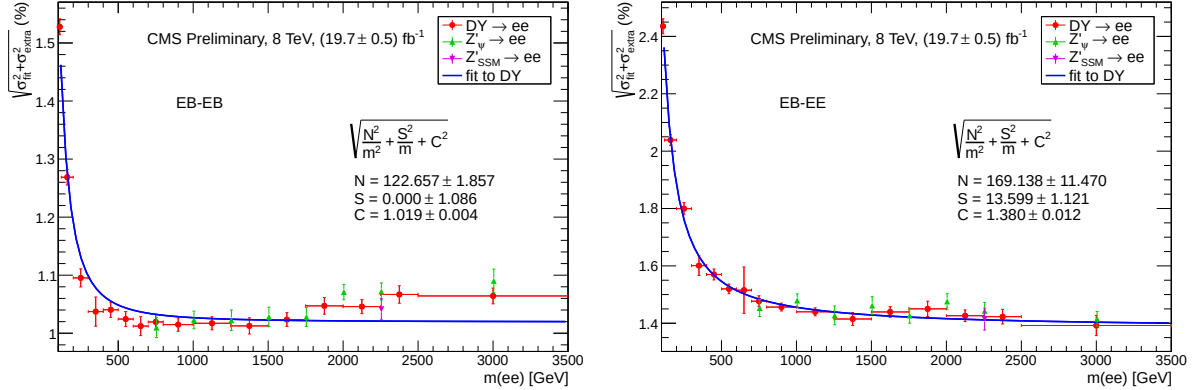


Figure 6.4: Evolution of the dielectron mass resolution as a function of the dielectron mass separately for pairs where both electrons are reconstructed in the ECAL barrel and where one of them is reconstructed in the endcaps. The figure compares results on a Drell-Yan simulated samples and on two signal simulated samples: the Z'_ψ and the Z'_{SSM} . A fit to the Drell-Yan points is performed using a function analogous to equation 2.6.

and a function of the same form is fitted. The results presented in figure 6.4 were obtained with Drell-Yan simulations. The results for two signal models, the Z'_ψ and the Z'_{SSM} are also shown for comparison. A good agreement is found. The resolution reaches asymptotically 1.2% and 1.8% at high mass for the barrel-barrel and barrel-endcaps categories. The mass resolution curve is an important input to the limit calculation. To take into account the difference between data and simulation measured at the Z peak, the extra term $\sigma_{DCB}(extra)$ is added quadratically to the functions showed in figure 6.4 and is conservatively assumed to be independent of the mass.

6.2.1 Saturation of the ECAL crystals

For very high energy electrons, the energy deposited in a single ECAL crystal can exceed the maximal value that the crystal can measure: the crystal saturates. As discussed in section 2.2.5.1, this happens for an energy deposit above 1.7 TeV in the barrel and 3.0 TeV in the endcaps. In the analysis of the 2012 data, no such electrons were found and the ECAL saturation has no impact on the results presented in this chapter.

A technique was nevertheless developed to measure the energy deposited in those crystals using the information from neighbours crystals [122, 123]. Starting from a 5×5 crystal matrix centered on the saturated crystal, one defines four 2×5 crystal dominos, $D_{\eta/x,l}$, $D_{\eta/x,h}$, $D_{\phi/y,l}$ and $D_{\phi/y,h}$ as illustrated in figure 6.5 (crystals are labelled by the η and ϕ coordinates in the barrel and by the x and y coordinates in the endcaps).

The two following variables are furthermore defined:

$$X = \ln \frac{\sum_{i \in D_{\eta/x,h}} E_i}{\sum_{i \in D_{\eta/x,l}} E_i} \quad (6.3)$$

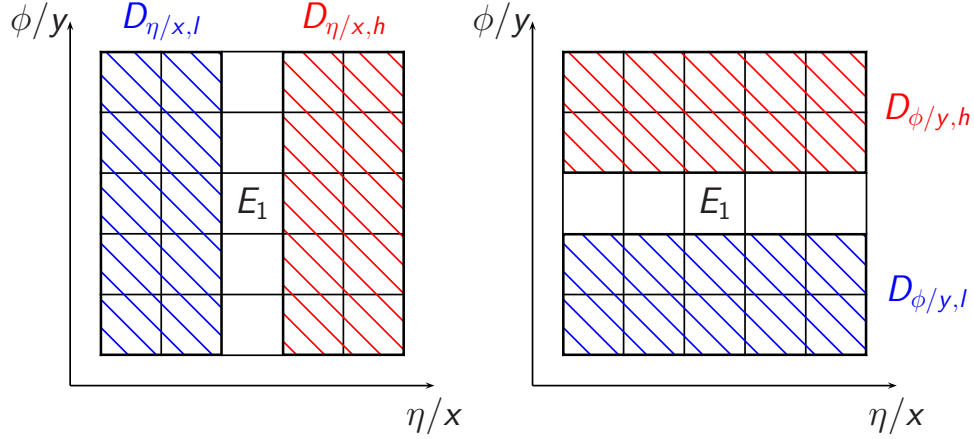


Figure 6.5: Definition of the four 2×5 crystal dominos used to estimate the energy of the central crystal, E_1 , in case the central crystal has a positive η or positive x and y . Crystals are labelled by the η and ϕ coordinates in the barrel and by the x and y coordinates in the endcaps. If η/x is negative, $D_{\eta/x,h}$ and $D_{\eta/x,l}$ are reverted. If y is negative, $D_{y,h}$ and $D_{y,l}$ are reverted.

$$Y = \ln \frac{\sum_{i \in D_{\phi/y,h}} E_i}{\sum_{i \in D_{\phi/y,l}} E_i} \quad (6.4)$$

where E_i is the energy of the crystal i . A parametric function $F(X, Y)$ is then fitted to the distribution of the ratio of the energy in the central crystal, E_1 , over the total energy in the 5×5 matrix, versus X and Y using electrons from simulations where the central crystal is not saturated, providing an estimate of the central crystal energy, denoted E_1^{rec} . The resolution of the method can be estimated from the distribution of $(E_1 - E_1^{rec})$.

Figure 6.6 presents the distribution of $(E_1 - E_1^{rec}) / \langle E_1 \rangle$ for electrons where the highest energy in any 2×5 crystal domino including the central crystal is above 500 GeV and below 2 TeV in the data and in simulation, where $\langle E_1 \rangle$ represents the mean E_1 value in the sample. The event selection requires two GSF candidates passing the high E_T electron selection.

The technique also provides a way to check the ECAL response to high energy signals: the fact that the data and simulation distributions show a fair agreement gives confidence that the ECAL response is well modeled by the simulation.

6.3 Background

The Standard Model processes leading to the reconstruction of dielectron pairs with high invariant mass can be split into three categories, which are, by decreasing order of importance, the annihilation of a $q\bar{q}$ pair (Drell-Yan process) or a $\gamma\gamma$ pair (“photon induced” process) into a dilepton pair, other processes involving two electrons in the final state and events where at least one of the reconstructed electrons is a fake.

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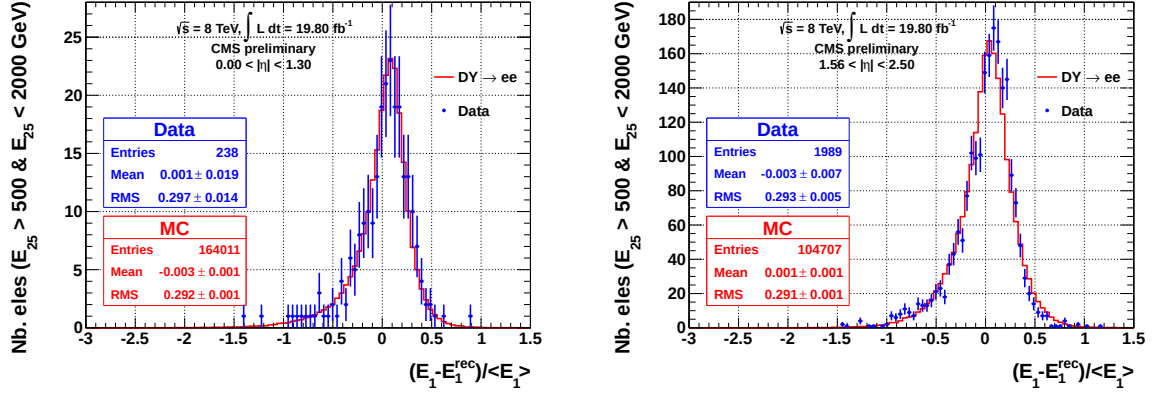


Figure 6.6: Distribution of $(E_1 - E_1^{rec}) / \langle E_1 \rangle$ for electrons where the highest energy in any 2×5 crystal domino including the central crystal is above 500 GeV and below 2 TeV, in the data (blue dots) and in a Drell-Yan simulated sample (red histogram). Results are presented separately for barrel (left) and endcaps (right) electrons.

6.3.1 Drell-Yan and photon induced processes

The Drell-Yan process represents the first contribution to the final dielectron mass spectrum and accounts for around 80 to 90% of the total events. Since the initial and final states can be identical to the signal studied, it is irreducible and interferences effects are even in theory possible. However, under the assumption of a narrow signal width, these effects are expected to be small and are neglected.

The Drell-Yan contribution is estimated from simulations using the Powheg NLO generator [96] for the hard process description, interfaced with Pythia for the description of the fragmentation and hadronization. The simulated events are reweighted using a global k-factor corresponding to the NNLO/NLO cross sections obtained by the FEWZ calculator [97] at the Z pole.

The Drell-Yan contribution is normalized to the data using events in the Z peak region ($60 < M_{e\bar{e}} < 120$ GeV). This normalization procedure is used to get rid of the uncertainty on the luminosity measurement and the limits calculated in section 6.5 will be set on the ratio of the cross section of a new heavy resonance over the cross section measured at the Z peak.

To check the accurate description of the data by the simulations, a measurement of the cross section of the Drell-Yan process in the Z peak region is performed, for the barrel-barrel and barrel-endcaps categories separately. The cross section, σ , is obtained as:

$$\sigma = \frac{N_{data} - N_{bkg}}{A \times \epsilon^2 \times SF_\epsilon^2 \times \epsilon_{trg} \times \int \mathcal{L} dt} \quad (6.5)$$

where N_{data} is the number of dielectrons pairs with $60 < M_{e\bar{e}} < 120$ GeV in the data, N_{bkg} is the expectation from all the Standard Model processes except the Drell-Yan process (as detailed in the next sections), $A \times \epsilon^2$ is the product of the acceptance times the electron efficiencies obtained from the simulation, SF_ϵ is the data over simulation correction factor

	barrel-barrel	barrel-endcaps
N_{data}	2138549	798481
N_{bkg}	10318 ± 689	3745 ± 300
$A \times \epsilon$	0.0973 ± 0.0010	0.0382
SF_ϵ	0.994 ± 0.012	0.976 ± 0.009
ϵ_{trg}	0.9930 ± 0.0004	0.9926 ± 0.0003
$\int \mathcal{L} dt \text{ (fb}^{-1}\text{)}$	19.7 ± 0.5	19.7 ± 0.5
$\sigma \text{ (pb)}$	$1125 \pm 15 \text{ (syst.)} \pm 29 \text{ (lumi)}$	$1090 \pm 15 \text{ (syst.)} \pm 28 \text{ (lumi)}$

Table 6.2: The various inputs to the measurement of the Drell-Yan cross section in the Z peak region ($60 < M_{e\bar{e}} < 120$ GeV) using pairs in the barrel-barrel (left) or barrel-endcaps (right) category. The last row shows the resulting cross section with its uncertainty. The uncertainty on the luminosity is shown apart from the other systematic uncertainties.

to the efficiency, obtained from table 5.3. ϵ_{trg} is the trigger efficiency, measured in data events and given at the end of section 5.4.2. and $\int \mathcal{L} dt$ is the integrated luminosity. The various inputs and the measured cross sections are shown in table 6.2.

The results are in good agreement with the theoretical cross section value of 1117 pb^{-1} , obtained with the FEWZ calculator at NNLO [97]. A 3% difference is however observed between the two categories. The origin of the discrepancy was found to be related to the electron energy scale discussed in section 6.2. In the endcaps, the electron energy measurement is slightly shifted up in the simulations. The effect on the acceptance of the E_T cut was measured to be 2%. The two measurements then agree within the uncertainties.

In addition to the Z peak, the beginning of the high mass tail of the Drell-Yan process, corresponding to invariant masses between 120 GeV and 200 GeV and where no new physics is expected is used as a control region as detailed in section 6.4.

The “photon induced” process is a process that has recently been the object of increasing interest [124]. The treatment of the quantum electrodynamics (QED) corrections to the parton distribution functions of quarks and gluons introduces a parton distribution function for the photon as well. Dilepton production through the annihilation of two photons at the LHC is therefore possible. It has been shown that this process is potentially sizeable at high mass (above 1 TeV) and could even become the dominant contribution from the Standard Model well above 2 TeV. Its theoretical uncertainty however is still large. In the current analysis, this process has not been simulated and it was assumed that its contribution is covered by the systematic uncertainty assigned to the total background contribution (*cf.* section 6.5.1.2). A better description of this process should be reached in the future. As discussed later on, the total background normalization has however little impact on a narrow resonance search such as the one performed in this thesis.

6.3.2 Other backgrounds with electrons in the final states

In addition to the Drell-Yan events, other processes lead to dielectron pairs in the final states. The main contributions for masses above the Z pole come from leptonic decays of

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$t\bar{t}$ or diboson (WW, WZ, ZZ) pairs with a small contribution from leptonic decays of a top or antitop quark produced in association with a W boson. In the Z pole region and below, a contribution from leptonic decays of ditau pairs is also expected.

In the final dielectron mass spectrum, all these processes are estimated from simulations using NLO or NNLO cross sections (*cf.* table 3.1). Since all of them involve a leptonic decay of a real or virtual W or Z boson, their normalization and shape can be checked in the data by considering the invariant mass spectrum of pairs of an electron and a muon ($e\text{-}\mu$ channel). Those processes are indeed dominant for this final state and their contribution, before any acceptance and efficiency effects are taken into account, is exactly twice their contribution in the dielectron channel.

Figure 6.7 compares the electron-muon mass spectrum observed in the data and in simulations. The event selection consists of a GSF electron candidate passing the high E_T selection and a well reconstructed muon with a transverse energy above 45 GeV with $|\eta| < 2.1$ passing the selection presented in table 2.2. The trigger used to record the data requires a muon and an electron or a photon object at the HLT. In addition to the processes already mentioned, depicted in red in figure 6.7, a small contribution from the Drell-Yan process ($\gamma/Z \rightarrow e\bar{e}/\mu\bar{\mu}$) where one of the lepton is misidentified as another lepton is expected at the Z pole (blue). Finally, W+jets (where the W decays into a muon and a neutrino and the jet fakes an electron) and multijets events also contribute a little (yellow). The former is estimated from simulations whereas the latter is estimated from the data as twice the difference, in the mass spectrum of $e\text{-}\mu$ pairs where the electron and the muon carries a same charge, between the data and the expected yields obtained from simulations for the other processes.

The $e\text{-}\mu$ invariant mass spectrum for same sign and opposite sign pairs are shown separately and together in figure 6.7. The ratio of the data over the Standard Model expectations is drawn on the bottom of the plots and a good agreement is observed in the whole spectrum.

6.3.3 Fake electrons background

The last contribution to the dielectron mass spectrum comes from events where at least one of the reconstructed electron is a fake. The processes involved are mainly W+jets, photon+jets and multijets. To measure their contribution, a data driven technique is used.

The starting point is the measurement of the probability for a jet to fake an electron. To do so, a data sample of events containing only one GSF electron candidate with $E_T > 10$ GeV is built using various triggers requiring one electron or photon candidate. This sample is mainly composed of multijets events and the jet fake rate is defined as the probability for a GSF electron candidate in this sample passing a loose preselection to pass the high E_T electron selection. The preselection consists of a series of cuts mainly emulating the HLT_DoubleEle33_CaloIdL_GsfTrkIdVL trigger. It is shown in table 6.3.

A small fraction of the events in the data sample comes from processes such as W+jets which can contain an electron in the final state. These events would bias the fake rate measurement and their contribution is removed using simulations. The fake rate is parametrized as a function of the GSF electron candidate E_T and η . It is shown as

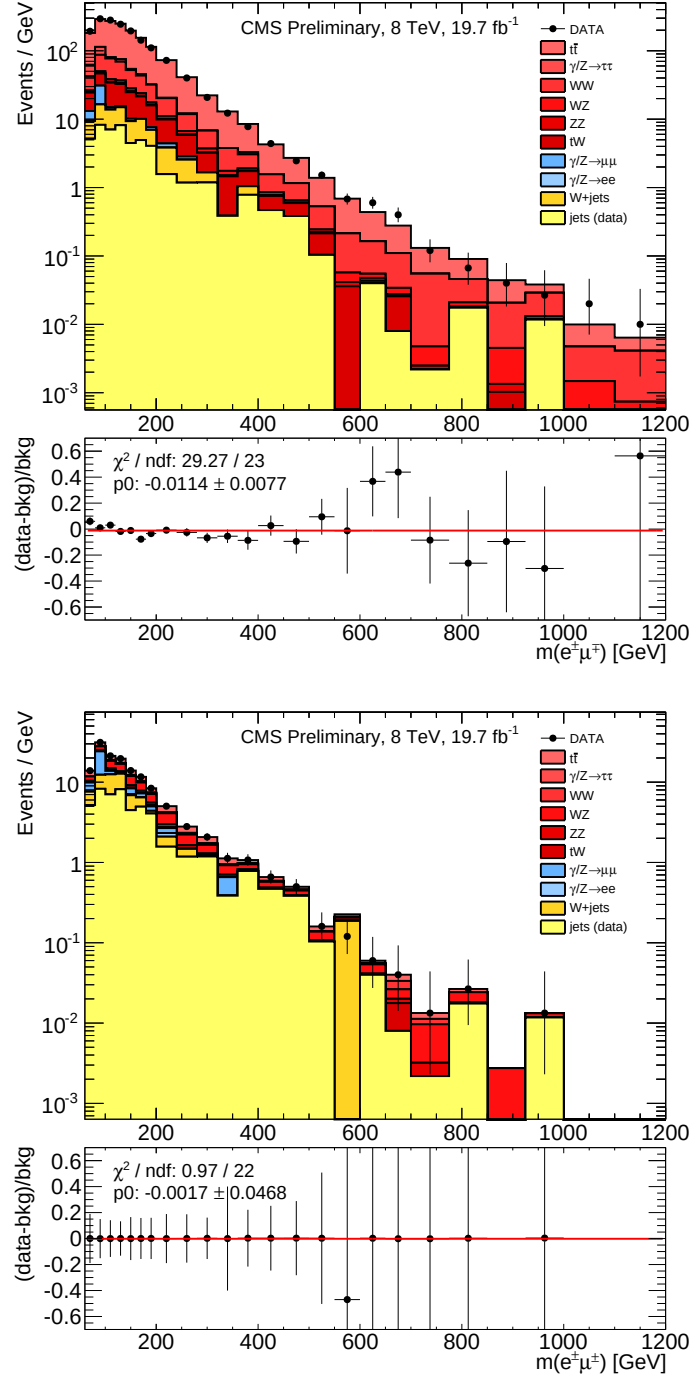


Figure 6.7: The invariant mass spectrum of electron-muon pairs, for pairs where the two leptons carry an opposite charge (top) or a same charge (bottom). The ratio $(data - bkg)/bkg$, where $data$ indicates the number of events observed in the data and bkg the expectation from the various Standard Model processes is given at the bottom of each plot. The contribution from multijets is estimated from the difference, in the same sign pair spectrum, between the data and the simulations of the other Standard Model processes.

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Variable	Barrel	Endcaps
H/E	<0.15	<0.1
$\sigma_{i\eta i\eta}$	<0.014	<0.035
Number of missing hits	≤ 1	≤ 1
d_{xy}	< 0.02	< 0.05

Table 6.3: The cuts entering the preselection applied to GSF electron candidates for the measurement of the jet fake rate. All the variables are defined in section 4.3.

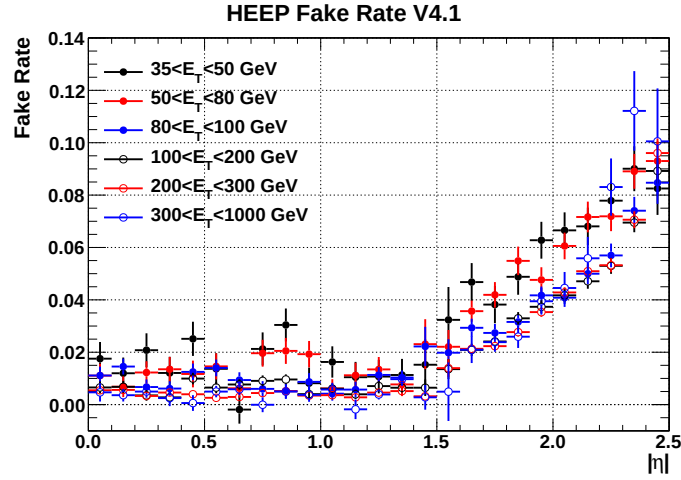


Figure 6.8: The jet fake rate measured as a function of the GSF electron candidate $|\eta|$ for various ranges in E_T .

a function of η for various E_T ranges in figure 6.8. In the barrel, no η dependency is observed. In the endcaps, a linear behaviour is seen, the slope being independent of E_T . The E_T dependency is shown in figure 6.9 and a linear fit is performed in three E_T regions. A 2D parametrization is therefore performed which uses the fact that the evolution of the fake rate with η does not depend on E_T . The parametrization is shown in table 6.4.

The contribution to the final dielectron mass spectrum from events where a fake electron is present can be estimated using two different samples enriched in multijets (where both electrons are fakes) and “W+jets like” events (with only one of the electron being a fake)¹:

- Using the mass spectrum obtained from pairs of GSF electron candidates that do not pass the high E_T electron selection (“2 fail” sample), weighted by the product of the fake rates associated to each of the two electrons². This sample is completely dom-

¹The case of photon+jets also enters the “W+jets like” category: in those events, one of the GSF electron candidates is, in most of the cases, coming from the conversion of the photon in the first layers of the tracker detector, as explained in section 4.3.

² The exact weight is $\frac{\zeta_1}{1-\zeta_1} \times \frac{\zeta_2}{1-\zeta_2}$, where ζ_1 and ζ_2 are the fake rate associated to each GSF electron candidate, since fakes passing the selection are not included in the sample. A similar remark holds for the “1 pass 1 fail” sample below

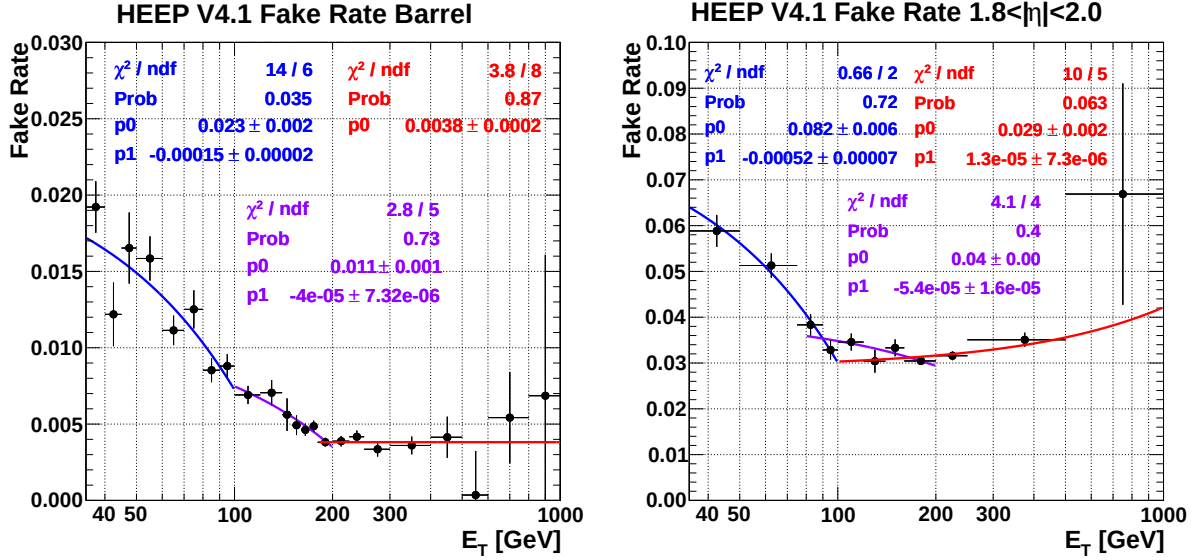


Figure 6.9: The jet fake rate measured as a function of the GSF electron candidate E_T for electrons reconstructed in the barrel (left) or endcaps (right). A fit to a straight line is performed for three E_T ranges.

Region	E_T range (GeV)	Parametrization
Barrel	$35 < E_T < 98.0$	$0.0226 - 0.000153 \times E_T$
	$98.0 < E_T < 191.9$	$0.0115 - 3.98 \times 10^{-5} \times E_T$
	$E_T > 191.9$	0.00382
Endcaps	$35 < E_T < 89.9$	$0.0823 - 0.000522 \times E_T + (\eta - 1.9) \times 0.065$
	$89.9 < E_T < 166.4$	$0.0403 - 5.45 \times 10^{-5} \times E_T + (\eta - 1.9) \times 0.065$
	$E_T > 166.4$	$0.0290 + 1.32 \times 10^{-5} \times E_T + (\eta - 1.9) \times 0.065$

Table 6.4: Parametrization of the jet fake rate as a function of the GSF electron candidate E_T and η .

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inated by multijets events and no background subtraction from processes involving an electron is needed. This procedure however does not estimate the contribution from W+jets or other processes where one of the two electrons is not a fake.

- Using the mass spectrum obtained from pairs of GSF electron candidates where exactly one of them passes the high E_T electron selection (“1 pass 1 fail” sample), weighted by the fake rate associated to the candidate that fails the selection. This procedure estimates both the multijets and W+jets like contribution. However a double counting happens for multijets since any of the candidate can be selected. In addition, this sample contains a sizeable contribution from Drell-Yan events which must be subtracted.

The technique used here mixes the two approaches: the final contribution is obtained from the “1 pass 1 fail” distribution subtracting the “2 fail” distribution to avoid double counting. The contribution from Drell-Yan events is estimated and subtracted using simulations.

To check the precision of the method, the invariant mass distribution of the “1 pass 1 fail” pairs in the data was compared to the predictions from the Standard Model processes, where the contribution from the multijets background is estimated by applying the fake rate on the “2 fail” distribution. The results are shown in figure 6.10 for the barrel-barrel and barrel-endcaps categories separately. Overall the data and expectations agree reasonably well as can be seen from the right plots which shows the data over simulation expectations ratio minus 1 for each bin in mass. Since the data and simulations agree within 40 %, a systematic uncertainty of 40 % is assigned to the jet fake rate measurement.

6.3.4 Total background

The shape of the invariant mass spectrum corresponding to the various Standard Model backgrounds described in the previous sections is one of the inputs to the limit setting procedure described in section 6.5. This total mass spectrum is shown in figure 6.11 and the results of a fit to a parametric function f of the form:

$$f(M) = e^{p_0 + p_1 \times M - p_2 \times M^2} M^{p_3} \quad (6.6)$$

are shown. The fit is performed for masses between 200 and 3500 GeV. The choice of the parametric function has no physical reason but is found to well model the expected mass distribution.

6.4 Invariant Mass spectrum

The invariant mass spectrum measured in the data is compared to the Standard Model expectations in figure 6.12 for the barrel-barrel and barrel-endcaps separately and together. The distributions in the simulations are normalized to the data using events at the Z peak as explained in section 6.3.1. The difference between the number of events

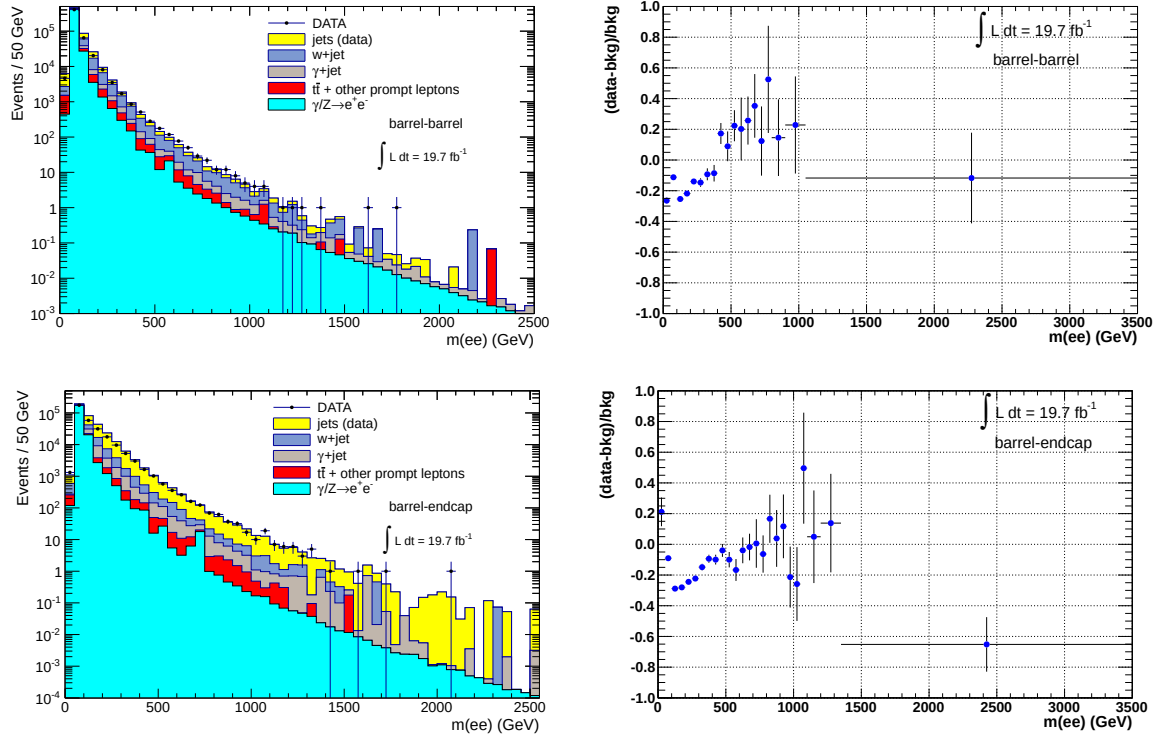


Figure 6.10: The invariant mass spectrum (left) of pairs of GSF electron candidates where one of them passes the high E_T electron selection the other not (black dots), separately for the barrel-barrel (top) and barrel-endcaps (bottom) categories. The plots on the right show the ratio of the number of events in the data minus the background expectation over the background expectation.

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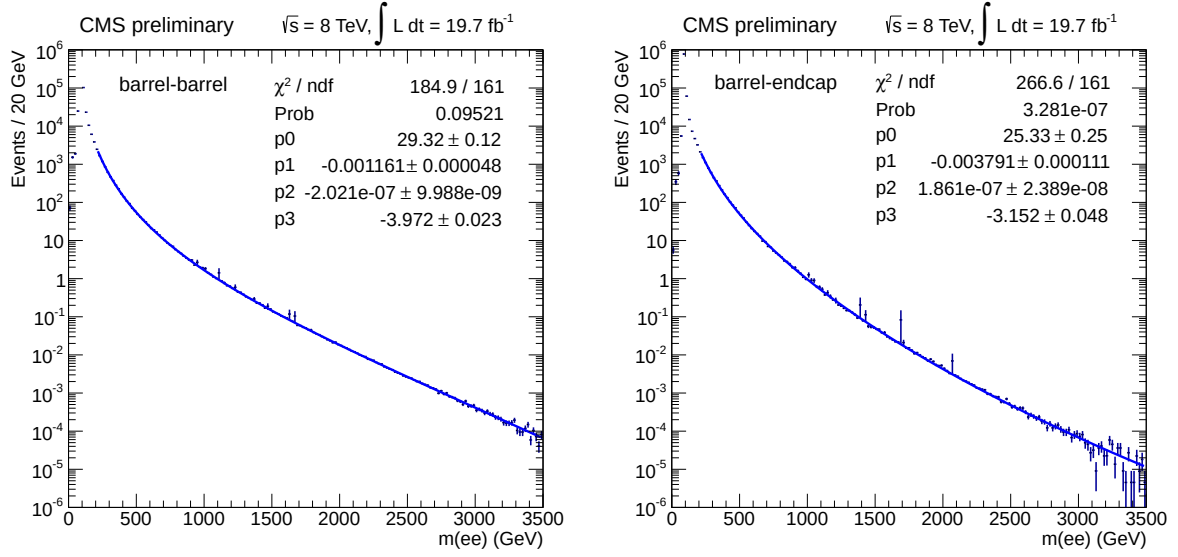


Figure 6.11: The expected contribution to the dielectron invariant mass spectrum from all the Standard Model processes separately for barrel-barrel and barrel-endcaps categories. A fit to a function given in equation 6.6 is also shown.

observed in the data and the expectations divided by the expectations is shown in figure 6.13 for dielectron pairs in the control region (masses between 120 and 200 GeV) and in the signal region (masses above 200 GeV). The Standard Model yields are calculated from the histograms in figure 6.12. Only statistical errors are drawn. No new resonant signal is observed. To quantify the global agreement between the mass spectrum in the data and in the simulations, a fit to a straight line is performed resulting in a data over simulation ratio compatible with 1 in both the control and signal regions.

An event display showing the highest mass dielectron pair recorded in the data is shown in figure 6.14. The reconstructed mass is 1790 GeV with a resolution of the order of 1%.

6.5 Statistical analysis and exclusion limits

Since no significant excess above the Standard Model predictions is seen, upper limits are set on the production cross section of new heavy dielectron resonances. The limits are calculated generically for narrow resonances, whose intrinsic width are of the same order of magnitude or smaller than the detector resolution. Since the acceptance and efficiencies differ for various spin hypotheses of the resonant signal, limits are computed for each of the cases, using the acceptance times efficiency parametrization from figure 6.2.

Those limits can reinterpreted as limits on the mass of resonances predicted by various benchmark models of new physics.

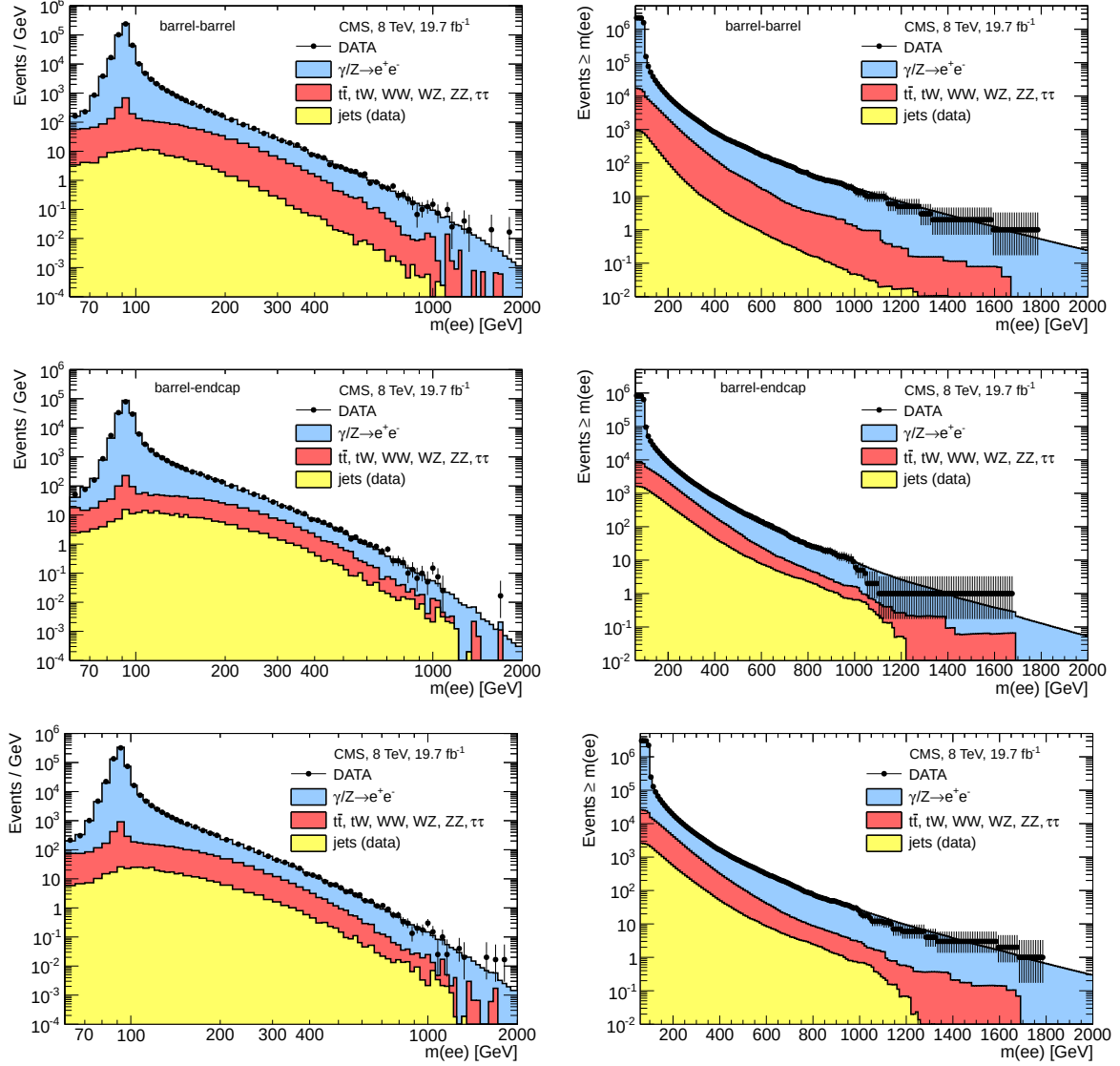


Figure 6.12: The dielectron invariant mass spectrum (left) in proton-proton collisions at $\sqrt{s} = 8$ TeV recorded by the CMS detector in 2012 (black dots) and the expected contributions from the Standard Model processes (histograms). Dielectron pairs where both electrons are reconstructed in the ECAL barrel (top) or where one of the electron is reconstructed in the endcaps (middle) are considered separately. The combination of the two categories is also shown (bottom). The right plots shows the associated cumulative distributions defined as the number of events above a given mass.

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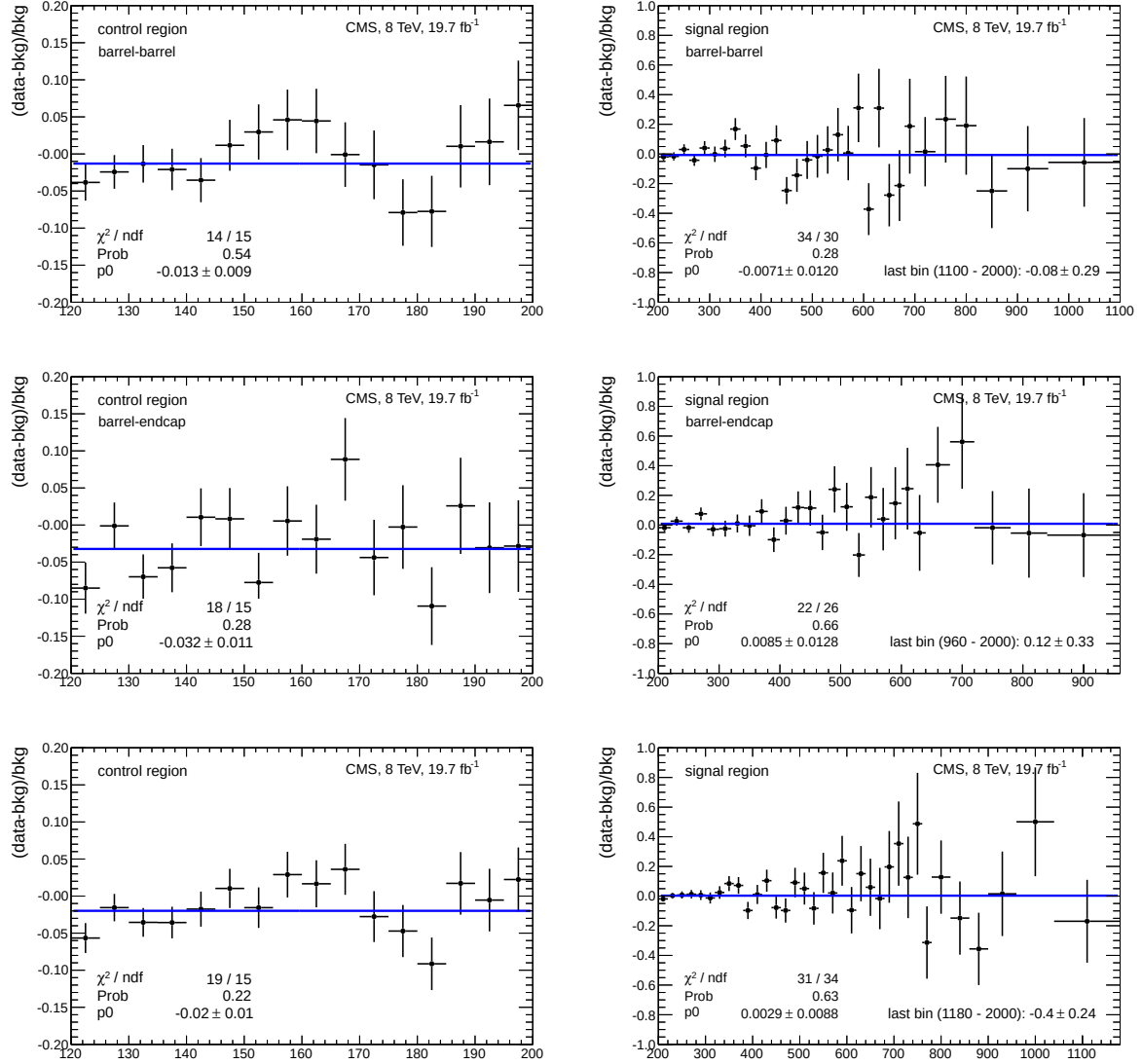


Figure 6.13: The observed number of data minus the expected background over the expected background for dielectrons pairs with an invariant mass between 120 and 200 GeV (“control region”, left) and above 200 GeV (“signal region”, right) for the barrel-barrel (top) and barrel-endcaps (middle) categories. A fit to a straight line is performed. The combination of the two categories is also shown (bottom).

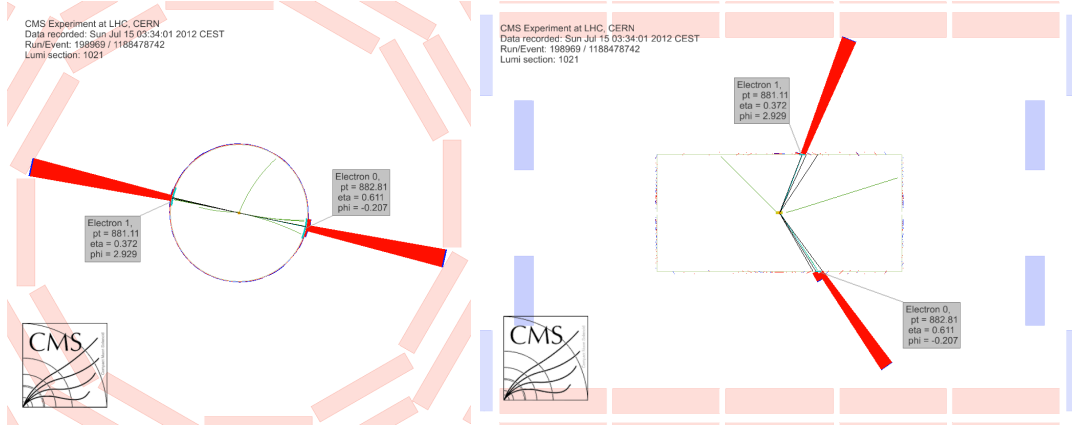


Figure 6.14: Transverse (left) and longitudinal (right) view of the highest mass dielectron pair recorded by CMS in proton-proton collisions at $\sqrt{s} = 8$ TeV. The mass of the pair is measured to be 1790 GeV. Only reconstructed tracks with a transverse impulsion above 3 GeV are shown.

6.5.1 Upper limits on the production cross section of new heavy dielectron resonances

A Bayesian approach [14, 125, 126] is followed to derive upper limits at the 95% confidence level on the production cross section of new heavy dielectron resonances. The limits are set using an unbinned likelihood based on the dielectron mass distribution. This allows one to exploit fully the very different shape of the signal (peak) and the background (falling continuum).

6.5.1.1 Procedure

One defines $f(m|\theta)$ as the probability density function (pdf) describing the dielectron mass (m) distribution. The pdf depends on various parameters of interest $\theta = \{\theta_1, \dots, \theta_L\}$ and on extra parameters $\nu = \{\nu_1, \dots, \nu_M\}$ subject to uncertainties. For a given set of N events, one denotes the corresponding set of the dielectron masses $\vec{m} = \{m_1, m_2, \dots, m_N\}$ and one defines the parametric likelihood as:

$$\mathcal{L}(\vec{m}|\theta, \nu) = \prod_{i=1}^N f(m_i|\theta, \nu) \quad (6.7)$$

The function $f(m_i|\theta, \nu)$ is the sum of the background shape $f_{bkg}(m_i|\theta, \nu)$ and the signal shape $f_{sig}(m_i|\theta, \nu)$, and the total event yield, μ , is the sum of the background yield, μ_{bkg} , and the signal yield, μ_{sig} . An extended likelihood is then defined as:

$$\mathcal{L}(\vec{m}|\theta, \nu) = \frac{\mu^N e^{-\mu}}{N!} \prod_{i=1}^N \left(\frac{\mu_{bkg}}{\mu} f_{bkg}(m_i|\theta, \nu) + \frac{\mu_{sig}}{\mu} f_{sig}(m_i|\theta, \nu) \right) \quad (6.8)$$

In this study the parameters θ consist of only one parameter: the signal yield. In order to get rid of various systematic uncertainties or to decrease their impact on the

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final results, another variable, R_σ , is rather used. It is defined as the ratio between the resonance (generically denominated Z') production cross section times branching ratio to the same quantity at the Z peak:

$$R_\sigma = \frac{\sigma_{pp \rightarrow Z' \rightarrow e\bar{e}}}{\sigma_{pp \rightarrow Z \rightarrow e\bar{e}}} \quad (6.9)$$

which is related to the signal yield through the relation:

$$\mu_{sig} = R_\sigma \frac{(A \times \epsilon)_{Z'}}{(A \times \epsilon)_Z} N_Z \quad (6.10)$$

where $\frac{(A \times \epsilon)_{Z'}}{(A \times \epsilon)_Z}$ is the ratio of the acceptance times efficiency between dielectron pairs produced at the Z' peak and at the Z peak, and N_Z is the number of events observed in the Z peak region.

Various parameters denoted ν in the likelihood are subject to an uncertainty. One example is the ratio $\frac{(A \times \epsilon)_{Z'}}{(A \times \epsilon)_Z}$. For a variable ν to which an uncertainty $\delta\nu$ is assigned, one introduces a nuisance parameter denoted β and the likelihood in equation 6.8 is modified in the following way:

$$\mathcal{L}(\vec{m}|\theta, \nu) \rightarrow \mathcal{L}(\vec{m}|\theta, \nu(1 + \delta\nu)^\beta) \times \frac{1}{\sqrt{2\pi}} e^{-\frac{\beta^2}{2}} \quad (6.11)$$

The variable ν is now floating but a Gaussian penalization factor is applied to the likelihood for any deviation with respect to its central value.

A Bayesian credible interval can then be built from this likelihood. Starting from a prior pdf $p(R_\sigma)$ for the R_σ parameter, which corresponds to the knowledge on the pdf before looking at the data, one can infer the posterior pdf $p(R_\sigma|\vec{m})$ using Bayes' Theorem:

$$p(R_\sigma|\vec{m})p(\vec{m}) = \mathcal{L}(\vec{m}|R_\sigma)p(R_\sigma) \quad (6.12)$$

where $p(R_\sigma|\vec{m})$, $\mathcal{L}(\vec{m}|R_\sigma)$ and $p(R_\sigma)$ are integrated over all the nuisance parameters. $p(R_\sigma|\vec{m})$ can therefore be expressed as:

$$p(R_\sigma|\vec{m}) = \frac{\mathcal{L}(\vec{m}|R_\sigma)p(R_\sigma)}{\int \mathcal{L}(\vec{m}|R_\sigma)p(R_\sigma)dR_\sigma} \quad (6.13)$$

$R_\sigma^{95\%}$, the 95% confidence level upper limit on R_σ , is the solution of:

$$\int_0^{R_\sigma^{95\%}} p(R_\sigma|\vec{m})dR_\sigma = 95\% \quad (6.14)$$

As mentioned earlier, the set of selected events is divided into two categories according to the region where the electrons are reconstructed (barrel-barrel and barrel-endcaps). Since the two categories show a different performance in terms of efficiency, mass resolution and background shape, a likelihood is defined for each category separately. The results can be combined by using the joint likelihood which is the product of the likelihood in each category.

	barrel-barrel	barrel-endcaps
N_Z	2138549	798481
$(A \times \epsilon)_Z$	0.0973	0.0382
$f_{bkg}(m)$	$e^{29.32-1.16 \times 10^{-3}m-2.02 \times 10^{-7}m^2}m^{-3.97}$	$e^{25.33-3.79 \times 10^{-3}m+1.86 \times 10^{-7}m^2}m^{-3.15}$
Mass resolution	$\sqrt{\frac{122.7^2}{m^2} + 1.019^2}$	$\sqrt{\frac{13.60^2}{m} + \frac{169.1^2}{m^2} + 1.38^2}$
$(A \times \epsilon)_{spin\ 0}$	$0.64 - \frac{2.7 \times 10^5}{m^2+6.6 \times 10^5}$	$2.4 - \frac{3.7 \times 10^4}{m+1.1 \times 10^4} + \frac{4.5 \times 10^6}{m^2+4.1 \times 10^6}$
$(A \times \epsilon)_{spin\ 1}$	$0.59 - \frac{2.9 \times 10^5}{m^2+7.5 \times 10^5}$	$0.06 - \frac{159}{m+345} + \frac{7.3 \times 10^5}{m^2+1.8 \times 10^6}$
$(A \times \epsilon)_{spin\ 2}$	$0.57 - \frac{3.0 \times 10^4}{m^2+1.3 \times 10^5}$	$-0.24 + \frac{1.2 \times 10^4}{m+3.5 \times 10^4} + \frac{6.6 \times 10^4}{m^2+7.5 \times 10^5}$

Table 6.5: The input parameters to the limit setting. All masses are in units of GeV.

The procedure that was just described allows one to compute the so called observed limits which are obtained using $\vec{m}$ measured in the data. Those limits are generally compared to the expected limits in the absence of a signal which are computed using pseudo-data generated according to the background only hypothesis. For each pseudo-experiment, a random value for β_{bkg} is drawn according to a normal centered distribution with a width given by $\delta\mu_{bkg}$. The total number of expected events n_{PE} is then estimated from a Poisson law with a mean equal to $\mu_{bkg} \cdot (1 + \delta\mu_{bkg})_{bkg}^\beta$ and the set $\vec{m} = \{m_1, m_2, \dots, m_{N_{PE}}\}$ is generated according to $f_{bkg}(m_i|\boldsymbol{\theta}, \boldsymbol{\nu})$. The procedure is repeated many times and the median expected limit and the ± 1 and $\pm 2 \sigma$ bands are estimated.

6.5.1.2 Inputs to the limit setting

The dielectron mass spectrum is scanned for masses between 300 and 3500 GeV. For a given resonance mass M_{res} , only events in the mass window $M_{res} \pm 6\sigma_G$ where σ_G is the detector resolution, are used in the likelihood, except if this leads to a number of events below 400. In that case, the lower mass window is adjusted so that at least 400 events are included. The choice of a narrow window ensures that, for high mass resonances, the background estimation is properly constrained in the signal region and is not completely driven by the first bins around 300 GeV where the statistics is much larger. The condition on the 400 events is imposed such that the statistical uncertainty on the normalization (5%) is small enough to have no impact on the obtained limits.

The background shape f_{bkg} in equation 6.8 is parametrized by the function given in equation 6.6. The signal shape f_{sig} is taken to be a convolution of a Breit-Wigner function, representing the intrinsic shape of the resonance and characterized by a width σ_{BW} , with a Gaussian distribution with a width σ_G modelling the experimental mass resolution. The value of σ_{BW} is model dependent. In the following calculation, a value of 0.6% of the resonance mass is assumed, which corresponds to the expectation for a Z'_ψ .

The various inputs to the limit setting are shown in table 6.5.

Three sources of systematic uncertainties are considered: the uncertainty of the evolution of the ratio $\frac{(A \times \epsilon)_{Z'}}{(A \times \epsilon)_Z}$, on the mass energy scale, and on the number of background events. Their values are shown in table 6.6. The uncertainty on the background contribution has several origins:

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Observable	Uncertainty (%)
$(A \times \epsilon)_{Z'}/(A \times \epsilon)_Z$	4 (barrel-barrel), 6 (barrel-endcaps)
Mass energy scale	1
Drell-Yan background (PDF)	$4.15 + 1.83 \times 10^{-3}.m + 2.68 \times 10^{-6}.m^2$
Drell-Yan background (NLO EW)	$-1 + 4.2 \times 10^{-3}.m$
$t\bar{t}$ and other prompt lepton background	2.6
Jet background	40

Table 6.6: The systematic uncertainties on input parameters for the limit setting. All masses are in units of GeV.

- The uncertainty on the number of Drell-Yan events in the signal region due to the uncertainty on the $\frac{(A \times \epsilon)_{Z'}}{(A \times \epsilon)_Z}$ ratio.
- The uncertainty on the Drell-Yan cross section due to the proton density functions (as discussed in section 3.1.1). Mass dependent uncertainties were derived using the PDF4LHC procedure [127] using various pdf sets (MSTW08, CTEQ12, CT10, NNPDF21) [128].
- The uncertainty on the Drell-Yan cross section due to the electroweak corrections at the next to leading order. A mass dependent ratio between the NLO and LO cross sections computed with the HORACE event generator [129] is taken as a systematic uncertainty.
- The uncertainty on the cross sections of the other prompt lepton backgrounds described in section 6.3.2 approximated by the $t\bar{t}$ cross section uncertainty [130] (2.6%).
- The uncertainty on the contribution from jets backgrounds estimated to be 40% in section 6.3.3.

Practically, because the Drell-Yan process represents the leading contribution to the invariant mass spectrum, the two last uncertainties turn out to be negligible. The full background uncertainty is therefore calculated as the quadratic sum of the three first contributions. It is extremely mass dependent: the background uncertainties at 500 GeV, 1 TeV and 3 TeV are respectively 7%, 10% and 35%. Despite these large values, the limits calculated here are not affected by the background uncertainty. This is due to the method that exploits the very different signal and background shapes: a broad excess or deficit of events with respect to the background expectation is still more compatible with the background hypothesis than with the signal hypothesis.

6.6 Lower limits on the mass of new resonances predicted by various models of new physics

The calculation of upper limits presented in the section 6.5.1 holds for any new physics model predicting a narrow resonance corresponding to the decay of a spin 0, 1 or 2 particle,

as long as the resonance can be considered narrow. For a given model where the resonance mass is free, it is then possible to convert the upper limit on the cross section to a (usually lower) limit on the resonance mass by finding the mass range where the theoretical cross section for the new resonance is above the observed limit.

The “narrow width condition” is crucial: the resonance must be described by the convolution of a Breit Wigner distribution and a Gaussian distribution. In the results presented below, an intrinsic width of 0.6% of the resonance mass, which corresponds to the Z'_ψ value, is assumed. It was checked that the limits are little affected for values up to 5%. This means that the contribution from the interference with the Drell-Yan process and from offshell production must remain negligible. For signals with a mass close to the kinematical limit of the proton parton distribution functions (*cf.* section 3.1.1), this is not always fully correct and the parton density functions can sensibly enhance the relative offshell production. An example is given by the Z'_{SSM} which has an intrinsic width of 3%. If one neglects the interference effects, 20% of the events have a mass more than 400 GeV below the resonant mass, for a Z'_{SSM} simulated with a mass of 2 TeV. The fraction increases to 63% for a mass of 3 TeV.

In order to take this effect into account, only the fraction of the production cross section corresponding to events in the resonance region should be taken into account. It was shown [131] that a window of $\pm 5\% \times \sqrt{s}$ around the resonant mass restores the narrow width approximation for a large fraction of models and only events in this mass window are therefore used when computing the theoretical cross sections referred to as “onshell cross section”.

The upper limits on the production cross sections are presented in figures 6.15, 6.16, and 6.17, for a spin 0, spin 1 and spin 2 resonances respectively. The theoretical onshell production cross sections for a Z'_{SSM} and Z'_ψ are also shown in figure 6.16 and for a Randall-Sundrum graviton with various couplings in figure 6.17. In the case of the Z'_{SSM} and Z'_ψ , the theoretical cross section obtained from Pythia are reweighted by a mass dependent K-factor that represents the ratio of the NNLO over LO cross section for Drell-Yan events estimated using the ZWPRODP calculator [132]. A constant K-factor of 1.3 is applied to the graviton cross sections.

A good agreement is seen between the observed limits and the median expectation. The fact that the oscillations of the observed limit around the median expectation cover the 95% interval of the expected limit gives confidence that the systematic uncertainties are well estimated. For masses above 2 TeV, the limit on R_σ reaches an asymptotic value of $2 \cdot 10^{-7}$ which corresponds to a signal cross section times branching ratio of around 0.2 fb. This happens in the region where no events are seen and where the acceptance times efficiency is flat and the expected background close to zero. This value can be verified through a simple calculation: in a background free region, the observation of 0 event excludes, at the 95% confidence level, a signal leading to more than 3 events¹. Given that the typical acceptance times efficiency in that region is 65% and the integrated luminosity is around 20 fb^{-1} , one indeed finds an upper cross section corresponding to 0.2 fb.

¹ A random variable distributed according to a Poisson distribution with a mean of 3 has $e^{-3}3^0/0! = 5\%$ of chances to take a null value.

6. SEARCH FOR NEW HEAVY DIELECTRON RESONANCES WITH THE DATA COLLECTED BY CMS AT $\sqrt{s} = 8$ TEV

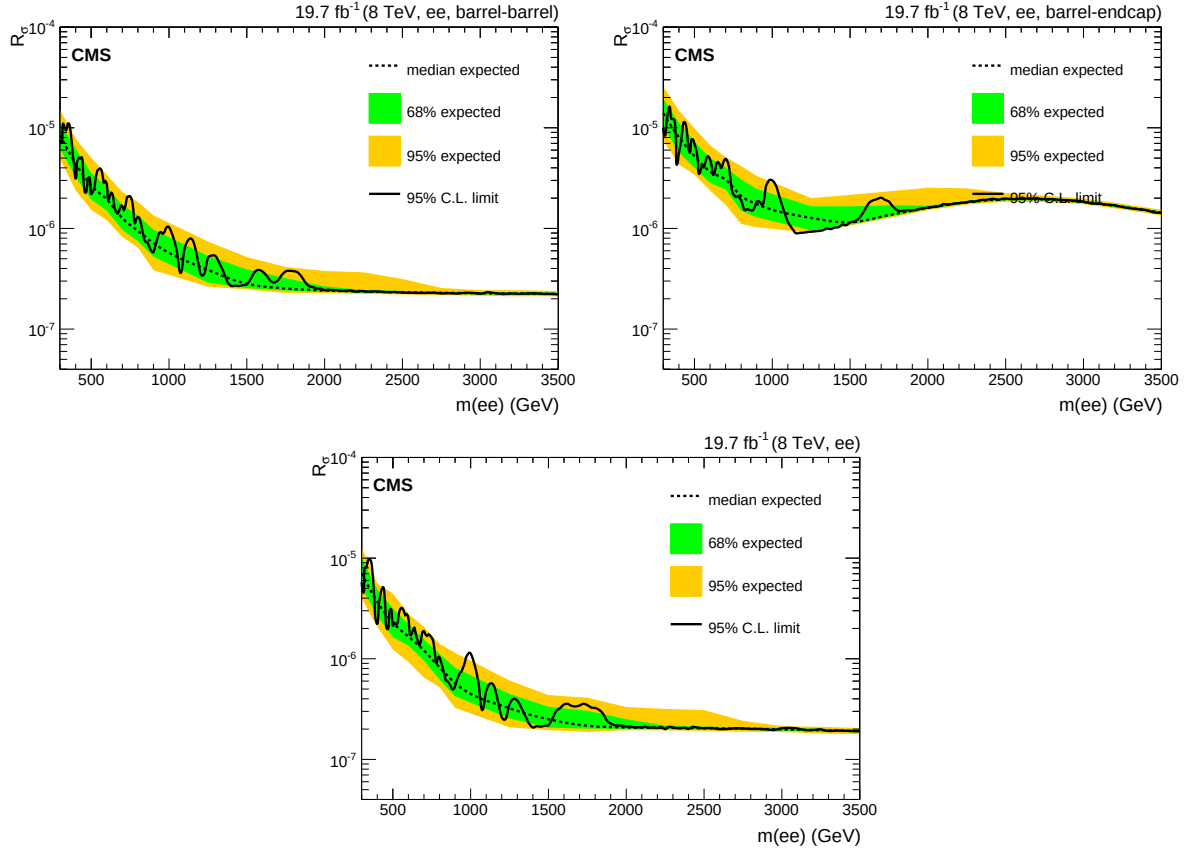


Figure 6.15: The 95% confidence level upper limit on R_σ for a spin 0 resonance coupling only to quarks as a function of the resonance mass using the full dataset recorded by CMS in 2012 (black). Also shown are the expected median limit (dashed line) and the centered 68% (green) and 95% (yellow) intervals for the expected limit. The results are shown separately for the barrel-barrel (top left) and barrel-endcaps (top right) categories and for the two categories combined (bottom).

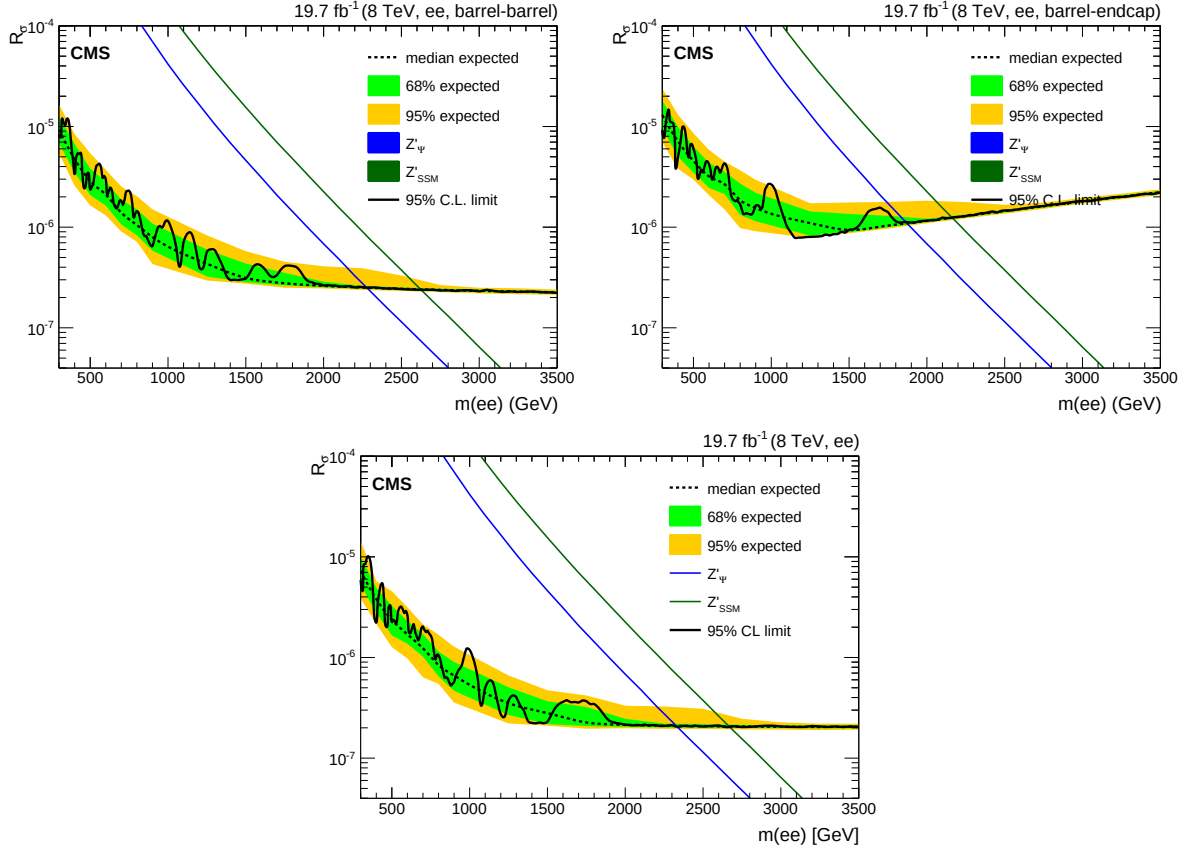


Figure 6.16: The 95% confidence level upper limit on R_σ for a spin 1 resonance as a function of the resonance mass using the full dataset recorded by CMS in 2012 (black). Also shown are the expected median limit (dashed line) and the centered 68% (green) and 95% (yellow) intervals for the expected limit. The blue (green) curve shows the production cross section times branching ratio of a Z'_ψ (Z'_{SSM}) boson as a function of its mass. The results are shown separately for the barrel-barrel (top left) and barrel-endcaps (top right) categories and for the two categories combined (bottom).

6. SEARCH FOR NEW HEAVY DIELECTRON RESONANCES WITH THE DATA COLLECTED BY CMS AT $\sqrt{s} = 8$ TEV

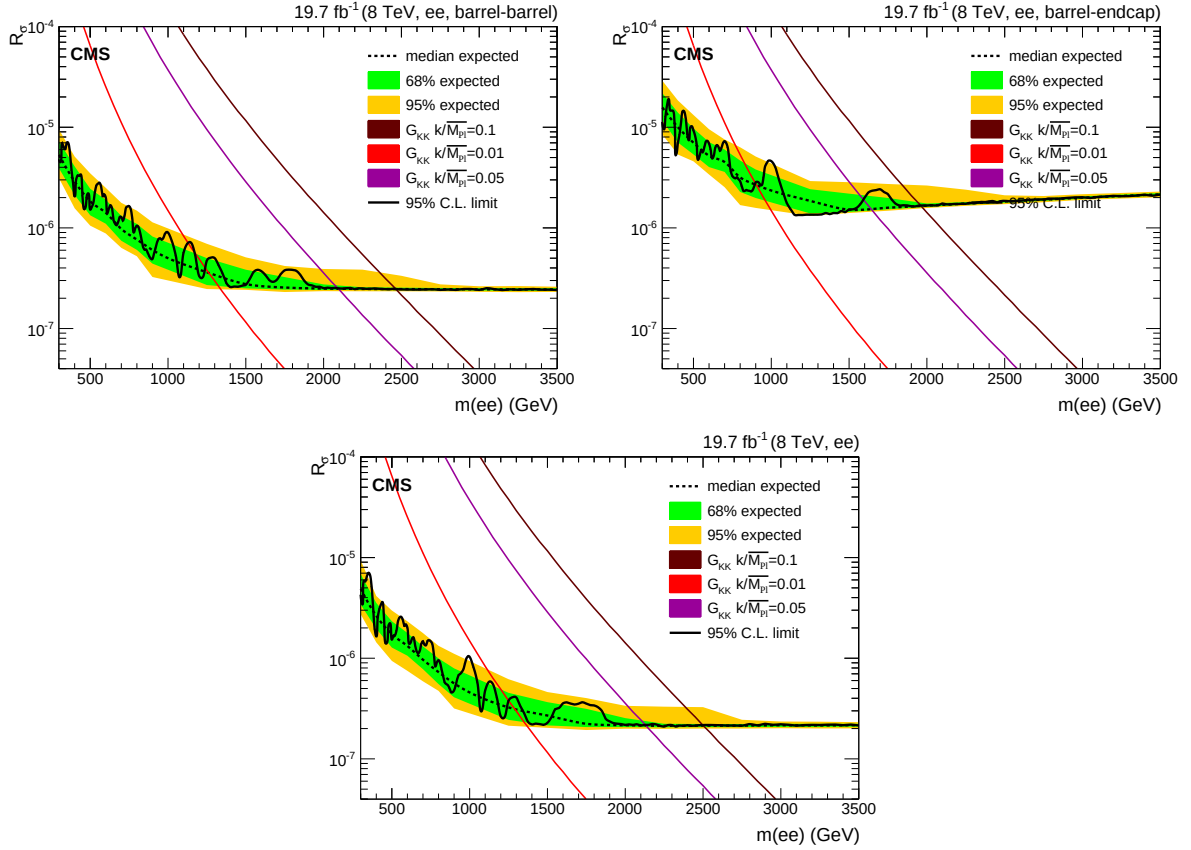


Figure 6.17: The 95% confidence level upper limit on R_σ for a spin 2 resonance with production mode similar to the Randall-Sundrum graviton as a function of the resonance mass using the full dataset recorded by CMS in 2012 (black). Also shown are the expected median limit (dashed line) and the centered 68% (green) and 95% (yellow) intervals for the expected limit. The blue (green) curve shows the production cross section times branching ratio of a Randall-Sundrum graviton with a coupling $c = 0.1$ ($c = 0.05$) as a function of its mass. The results are shown separately for the barrel-barrel (top left) and barrel-endcaps (top right) categories and for the two categories combined (bottom).

	barrel-barrel	barrel-endcaps	all dielectron	all dielectron + dimuon channel
Z'_ψ	2280	1870	2340	2570
Z'_{SSM}	2630	2170	2670	2900
$G_{RS}(c = 0.05)$	2090	1590	2130	2350
$G_{RS}(c = 0.1)$	2470	1960	2500	2730

Table 6.7: The lower limits on the mass of the various benchmark models considered for the barrel-barrel and barrel-endcaps categories separately and together. The last column presents the excluded mass obtained when combining the dielectron and dimuon channels.

6.6.0.3 Combination with the dimuon channel

The results of the search for heavy dielectron resonances are often combined with the dimuon channel assuming equal branching ratios between the two leptonic decays of the new resonance [4, 7]. The two channels have indeed similar sensitivities. The dimuon channel benefits from a better acceptance times efficiency and is less affected by the selection of fake muons but has a much worse mass resolution. Table 6.7 shows, for the various models considered, the lower limit at 95% confidence level on the resonance mass, for the dielectron analysis alone and after its combination with the dimuon channel.

6.7 Summary

This chapter details the search for a new heavy dielectron resonance using proton-proton collisions at $\sqrt{s} = 8$ TeV recorded by the CMS detector. In order to remain as model independent as possible and to maintain a high acceptance times efficiency, the event selection is very generic and requires only the presence of two electrons.

Except for the electron selection and its efficiency measurement which were described in details in chapters 4 and 5, all the aspects of the analysis were discussed. This includes the measurement of the dielectron mass resolution and the estimation of the contribution of the backgrounds from the Standard Model. Those studies are mainly performed with simulated samples but whenever possible, data driven techniques are used to correct or to check the simulations. Since the reconstruction and selection of electrons in the ECAL barrel and endcaps show a different performance in terms of efficiency, fake rejection and energy resolution, the dielectron pairs considered here are split into a barrel-barrel and a barrel-endcaps category.

The final dielectron mass spectrum measured in the data was scrutinized from the Z peak region (90 GeV) to masses up to several TeV. No significant deviation with respect to the background expectation was seen and upper limits were set on the ratio of the production cross section times branching ratio of a new heavy resonance to the corresponding value at the Z resonance. The limit setting procedure uses an unbinned likelihood based on the dielectron invariant mass distribution and exploits the difference in shape between the signal and the background. Limits are set for three different acceptance times efficiency scenarios corresponding to a spin 0, a spin 1 or a spin 2 resonance. For a given acceptance times efficiency, they are generic and apply to any signal leading to a local excess of events in the dielectron mass spectrum. The combination of the two electron categories allows one to exclude at the 95% confidence level a production cross section times branching ratio above 0.2 fb for resonances with a mass above 2 TeV.

The upper limits can be reinterpreted as a lower mass condition on resonances predicted by various new physics models. A Z'_ψ or a Z'_{SSM} below 2.34 TeV or 2.67 TeV respectively are excluded. Similarly, for the Randall-Sundrum graviton with a coupling $c=0.05$ ($c=0.1$), a lower limit of 2.13 TeV (2.50 TeV) is set.

The results presented here are usually combined with the analysis of the dimuon mass spectrum. Adding the dimuon channel brings the previous limits up to 2.57 TeV (Z'_ψ), 2.90 TeV (Z'_{SSM}), 2.35 TeV (G_{RS} , $c=0.05$) and 2.73 TeV (G_{RS} , $c=0.1$). Similar searches were conducted by the ATLAS Collaboration leading to comparable results [121].

Chapter 7

Study of the characterisation of a potential new dielectron resonance signal that could be discovered in the early data at $\sqrt{s} = 13$ and 14 TeV

In chapter 6, results for the search for new heavy narrow resonances decaying into a dielectron pair were presented using proton-proton collisions at $\sqrt{s} = 8$ TeV with the CMS detector. No significant excess above the Standard Model predictions was found and upper limits were set on the production cross sections of such phenomena. For the models of new massive resonances production described in section 1.4, those results can be turned into a lower limit on the mass of the resonance, reaching values up to 3 TeV in some cases. Similar results were obtained by the ATLAS collaboration. As of today, those limits are the world most stringent limits.

However, the foreseen increase of the LHC beam energy from 4 TeV to 6.5 or 7 TeV (*cf.* section 2.1.4) in the upcoming years will rapidly allow one to probe much higher mass values. To produce a particle with a mass in the TeV range, the fractions of the proton momentum carried by the two incoming partons, x_1 and x_2 , must be large and close to kinematical limit. In that region, the parton probability functions strongly vary with x_1 and x_2 . For two partons forming an invariant mass in the TeV range, increasing the beam energy by $\approx 60\text{-}75\%$ reduces the average x_1 and x_2 by the same amount, and will lead to a sizeable change in the cross section as explained in section 3.1.1. The ratio of the production cross section predicted by the Pythia generator for dielectron pairs from the Drell-Yan process with an invariant mass above 4 TeV at $\sqrt{s} = 13$ TeV and $\sqrt{s} = 8$ TeV is for instance greater than 100.

The expected minimal production cross section of a heavy dielectron resonance for a 5σ discovery are shown in figure 7.1 for various integrated luminosities and for proton-proton collisions at $\sqrt{s} = 14$ TeV [133]. A 5σ discovery is reached for a given signal model if:

$$p = \int_{N_{bkg}+N_{sig}}^{\infty} \text{Poisson}(x|N_{bkg})dx < 3.10^{-7} \quad (7.1)$$

where N_{sig} and N_{bkg} are respectively the number of signal and background events counted

7. STUDY OF THE CHARACTERISATION OF A POTENTIAL NEW DIELECTRON RESONANCE SIGNAL THAT COULD BE DISCOVERED IN THE EARLY DATA AT $\sqrt{s} = 13$ AND 14 TEV

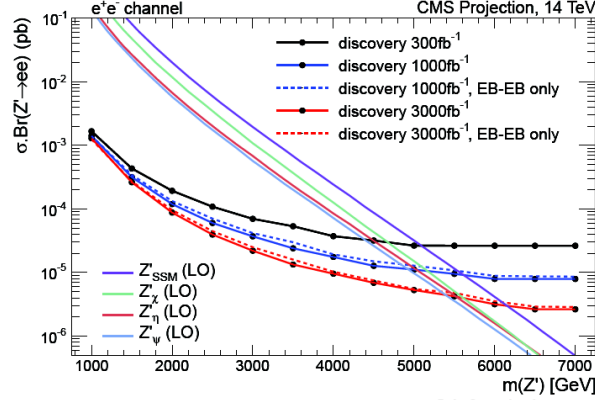


Figure 7.1: The minimal production cross section times branching ratio of a new heavy resonance decaying into a dielectron pair for a 5σ discovery, for three integrated luminosities for proton-proton collisions at $\sqrt{s} = 14$ TeV. The results are presented for a selection analogous to chapter 6 (full line) and with the extra condition that both electrons are reconstructed in the barrel (dashed line). The cross sections corresponding to some of the Z' models discussed in section 1.4.2.2 are also shown as a function of the Z' mass. Figure is taken from [133].

in a mass window containing 95% of the events in the signal peak, and if at least 5 signal events are found in that window. Figure 7.1 presents results for a selection analogous to the one used in chapter 6 (full line). A scenario where the ECAL endcaps are too affected by radiation damage to be used is also considered: in that case the two electrons are required to be reconstructed in the ECAL barrel (dashed line). The theoretical cross section obtained from Pythia for various Z' models are also indicated as a function of the mass of the Z' .

These projected limits are much stronger than the current ones and the quick emergence of a new signal, as soon as a few months after the LHC resumes in 2015 at $\sqrt{s} = 13$ TeV, is therefore possible.

If such a signal is to be found in the early data, the large integrated luminosity collected in the following years will allow one to measure its properties (mass, width, spin,...) in detail. This chapter studies the sensitivity that could be reached for its spin determination, the measurement of its forward-backward asymmetry and the measurement of its $gg-q\bar{q}$ production mode fractions.

Three luminosity scenarios are considered, corresponding to the total integrated luminosity expected to be collected at the end of the 2 upcoming runs of LHC (runs 2 and 3) and by the end of the LHC phase 2 (*cf.* table 2.1):

- 100 fb $^{-1}$ and $\sqrt{s} = 13$ TeV (run 2)
- 300 fb $^{-1}$ and $\sqrt{s} = 14$ TeV (run 3)
- 3000 fb $^{-1}$ and $\sqrt{s} = 14$ TeV (phase 2)

The measurements are based on the analysis of the lepton angular distributions and the dielectron pair rapidity. They are designed to be as model independent as possible and are applied to a couple of benchmark models to quantify the expected level of sensitivity that could be reached.

This chapter is structured as followed. Section 7.1 describes the various simulated samples used in the present study and details the event selection. Section 7.2 discusses the kinematics and the properties of $parton\ parton \rightarrow l\bar{l}$ processes and introduces the angular variables used in measurements presented in section 7.3.

7.1 Simulated samples and event selection

7.1.1 Simulated samples

The samples studied in this chapter were simulated with the Pythia generator (version 8) [94] for proton-proton collisions at $\sqrt{s} = 13$ TeV.

The only Standard Model process studied in this chapter is the Drell-Yan process, which represents the first contribution to high mass dilepton spectrum as observed in chapter 6. Two Drell-Yan (DY) samples were simulated with $M > 3$ and 4 TeV respectively.

The various signal samples considered are the Z'_{SSM} and the Z'_ψ described in section 1.4.2.2 and the Randall-Sundrum graviton G_{RS} described in section 1.4.3.3. The coupling parameter ($c = k/M_{Pl}$) for the G_{RS} is taken to be 0.1. All the samples were generated at a mass of 3 or 4 TeV without interference with the Drell-Yan process. These mass and couplings values are chosen to correspond to a signal not already excluded and whose yields are strong enough to provide enough statistics for the study performed here.

All the events were reconstructed with the full simulation of the CMS detector using the version 623 of the CMS software. Although the average number of pile-up interactions in table 2.1 is much larger in phase 2 than in Run 1, the same pile-up scenario (PU-S10) as for the simulation at $\sqrt{s} = 8$ TeV described in table 3.2 is assumed. The effect of the pile-up on the reconstruction and selection of electrons with energy above 1 TeV is however negligible.

All the distributions used in this chapter were obtained from the simulated samples at $\sqrt{s} = 13$ TeV. For the projections at $\sqrt{s} = 14$ TeV, the distributions were rescaled by the ratio of the 14 over 13 TeV cross sections.

The complete list of processes considered is indicated in table 7.1 with their cross sections at $\sqrt{s} = 13$ TeV and $\sqrt{s} = 14$ TeV.

7.1.2 Event selection

The event selection is similar to the one presented in chapter 6 and consists of two high E_T (> 35 GeV) GSF electron candidates passing the high E_T electron selection. Because of the sizeable mismeasurement on the charge for electrons with $E_T \approx 1$ TeV, an opposite charge of the two electrons is not required, except for the measurement of the forward-backward asymmetry which is only defined when the negative and positive lepton can be distinguished. The event must in addition fire the trigger described in table 4.2.

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Process	LO cross section (pb) $\sqrt{s} = 13$ TeV	LO cross section (pb) $\sqrt{s} = 14$ TeV
G_{RS} M=3 TeV	$6.1 \cdot 10^{-4}$	$8.7 \cdot 10^{-4}$
G_{RS} M=4 TeV	$5.5 \cdot 10^{-5}$	$9.3 \cdot 10^{-5}$
Z'_{SSM} M=3 TeV	$1.7 \cdot 10^{-3}$	$2.3 \cdot 10^{-3}$
Z'_{SSM} M=4 TeV	$2.6 \cdot 10^{-4}$	$3.6 \cdot 10^{-4}$
Z'_ψ M=3 TeV	$4.5 \cdot 10^{-4}$	$6.0 \cdot 10^{-4}$
Z'_ψ M=4 TeV	$5.1 \cdot 10^{-5}$	$7.8 \cdot 10^{-5}$
DY M>3 TeV	$1.2 \cdot 10^{-5}$	$1.8 \cdot 10^{-5}$
DY M>4 TeV	$1.2 \cdot 10^{-6}$	$1.9 \cdot 10^{-6}$

Table 7.1: List of the processes studied in this chapter, together with their cross sections computed at leading order with the Pythia generator for $\sqrt{s}=14$ TeV and $\sqrt{s}=13$ TeV.

The total acceptance times efficiency for the various signal models and the Drell-Yan process is shown in table 7.2, together with the expected yields for 100 fb^{-1} at $\sqrt{s} = 13$ TeV. For the signal samples, an additional cut is applied: only events for which $|M_{reco} - M_{res}| < 0.1 M_{res}$ are selected, where M_{reco} is the reconstructed mass and M_{res} the resonance mass. This defines a region where the background contamination and the possible interference effects with the Drell Yan process are low. For proton-proton collisions at 13 TeV, the production of 3 or 4 TeV resonances happens close to the kinematical limit of the pdfs. When their intrinsic width is large, off-shell production becomes sizeable and a large fraction of events are produced below the resonance peak. This is the case for the Z'_{SSM} for which the $|M_{reco} - M_{res}|$ cut removes an important fraction of the events. The Drell-Yan contribution in the signal region is estimated to be always below 3% of the signal yield. The hypothesis of a background free region is therefore made in the rest of this study.

7.2 Kinematics and properties of parton parton $\rightarrow \ell\bar{\ell}$ processes

7.2.1 Generalities

The annihilation into a dilepton pair of two partons with different energy but colinear can be described by three observables¹: the lepton pair invariant mass $M_{\ell\bar{\ell}}$, its rapidity $y_{\ell\bar{\ell}}$ defined as

$$y_{\ell\bar{\ell}} = \ln \frac{E_{\ell\bar{\ell}} - p_{z,\ell\bar{\ell}}}{E_{\ell\bar{\ell}} + p_{z,\ell\bar{\ell}}} \quad (7.2)$$

and the angle θ in the center of mass of the dilepton pair, between the negative lepton and one of the partons, randomly selected.

¹ Assuming the lepton and partons are massless.

Process	Acc. $\times$ Eff.	Yields	Yields (Opposite sign)
		$\sqrt{s} = 13 \text{ TeV}, 100 \text{ fb}^{-1}$	$\sqrt{s} = 13 \text{ TeV}, 100 \text{ fb}^{-1}$
$G_{RS} \text{ M}=3 \text{ TeV}$	0.644	39.3	36.8
$G_{RS} \text{ M}=4 \text{ TeV}$	0.637	3.5	3.2
$Z'_{SSM} \text{ M}=3 \text{ TeV}$	0.495	84.2	78.3
$Z'_{SSM} \text{ M}=4 \text{ TeV}$	0.372	9.7	8.9
$Z'_\psi \text{ M}=3 \text{ TeV}$	0.630	28.4	26.4
$Z'_\psi \text{ M}=4 \text{ TeV}$	0.596	3.0	2.8
DY $\text{M}>3 \text{ TeV}$	0.666	0.8	0.8
DY $\text{M}>4 \text{ TeV}$	0.673	0.1	0.1

Table 7.2: Total acceptance times efficiency and expected yields for 100 fb^{-1} of data collected at $\sqrt{s} = 13 \text{ TeV}$ for the different processes studied. For the signal samples, only events with a reconstructed mass compatible with the resonance mass within 10% are selected. The last column shows the yields when the two electrons are in addition required to carry an opposite charge.

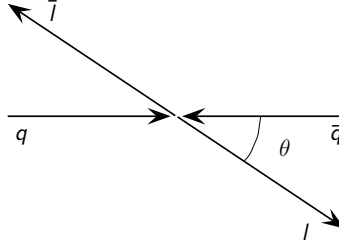


Figure 7.2: In a $q\bar{q} \rightarrow l\bar{l}$ process, θ is defined as the angle, in the center of mass of the dilepton pair, between the directions of the negative lepton and the quark.

In the Standard Model, the only existing process at tree level is the Drell-Yan process ($q\bar{q} \rightarrow \gamma/Z \rightarrow l\bar{l}$). In such a case of a $q\bar{q}$ annihilation, the selected parton is in addition requested to be the quark (*cf.* figure 7.2). The knowledge of the quark direction indeed allows one to study a possible preference of the lepton for this direction compared to the opposite. This is quantified by the forward-backward asymmetry A_{FB} already introduced in section 1.2.2 and given by:

$$A_{FB} = \frac{\sigma_F - \sigma_B}{\sigma_F + \sigma_B} \quad (7.3)$$

where $\sigma_F = \sigma_{\cos\theta>0}$ ($\sigma_B = \sigma_{\cos\theta<0}$) is the cross section corresponding to events where the negative lepton is emitted at an angle smaller (greater) than $\pi/2$ with respect to the quark direction. As partons are confined in hadrons in nature, one experimentally studies hadrons collisions. Only the incoming hadrons and outgoing leptons momenta are therefore measured. As a consequence, the exact measurement of θ (or $\cos\theta$) is not possible if at least one of the two incoming partons has a non zero transverse momentum. Moreover, the quark direction in the case of a $q\bar{q}$ annihilation is not directly known.

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7.2.2 The Collins-Soper frame

7.2.2.1 Original definition

The Collins-Soper (CS) frame [134] is the frame generally used for the study of the angular distribution of lepton pairs produced in the collision of two hadrons. It attempts to minimize the distortion of the angular distributions measured in the dilepton center of mass frame due to the transverse momenta of the two incoming partons. Let $\vec{p}_A$ and $\vec{p}_B$ be the momenta of the two colliding hadrons in the laboratory frame and let $\vec{p}'_A$ and $\vec{p}'_B$ be their values in a frame in the center of mass of the dilepton pair. The CS frame is defined as the frame in the center of mass of the dilepton pair for which:

- the z -axis bisects $\vec{p}'_A$ and $-\vec{p}'_B$.
- the x -axis is orthogonal to the z -axis, lies also in the plane spanned by $\vec{p}'_A$ and $\vec{p}'_B$ and its direction is chosen such that $p'_{x,A}$ and $p'_{x,B}$ are negative.

The construction of the frame is illustrated in figure 7.3. In this frame, one defines θ_{CS} as the angle between the negative lepton momentum and the z -axis. One also defines ϕ_{CS} as the azimuthal angle of the negative lepton momentum in the $x - y$ plane. One can show that θ_{CS} can be expressed in term of the lepton variables in the laboratory frame as¹:

$$\cos \theta_{CS,raw} = 2 \frac{E_l \cdot p_{z,\bar{l}} - E_{\bar{l}} \cdot p_{z,l}}{M_{l\bar{l}} \sqrt{M_{l\bar{l}}^2 + P_{T,l\bar{l}}^2}} \quad (7.4)$$

When the transverse boost of the lepton pair is small compared to its longitudinal boost, θ_{CS} provides a good approximation of θ , the angle between the negative lepton and one the incoming partons in their center of mass. At this level, no distinction is made between the two incoming partons.

7.2.2.2 Adding the quark direction

For $q\bar{q}$ annihilation one is interested in measuring θ_{CS} with respect to the quark direction. This is achieved by requiring the hadron A to be the hadron that provides the quark in the collision, or equivalently by using:

$$\cos \theta_{CS} = \frac{p_{z,q}}{|p_{z,q}|} 2 \frac{E_l \cdot p_{z,\bar{l}} - E_{\bar{l}} \cdot p_{z,l}}{M_{l\bar{l}} \sqrt{M_{l\bar{l}}^2 + P_{T,l\bar{l}}^2}} \quad (7.5)$$

Experimentally, the quark information cannot directly be accessed but in symmetric proton-proton collisions, the ambiguity of the direction of the quark and the antiquark can be partially removed. On average, a quark carries a bigger fraction of the proton energy than an antiquark and the dilepton pair is boosted in that direction. One therefore requires that proton A is the proton whose direction is the same as the longitudinal boost of the dilepton pair in the laboratory frame. This is equivalent to the following definition:

$$\cos \theta_{CS_{meas}} = \frac{p_{z,l\bar{l}}}{|p_{z,l\bar{l}}|} 2 \frac{E_l \cdot p_{z,\bar{l}} - E_{\bar{l}} \cdot p_{z,l}}{M_{l\bar{l}} \sqrt{M_{l\bar{l}}^2 + P_{T,l\bar{l}}^2}} \quad (7.6)$$

¹ The demonstration is given in appendix D.

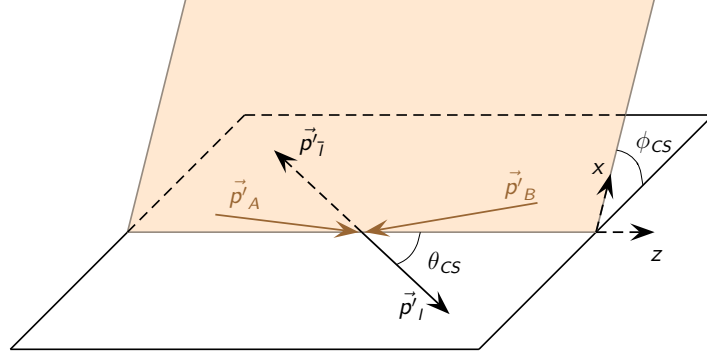


Figure 7.3: Definition of the Collins-Soper frame and of θ_{CS} and ϕ_{CS} . The frame lies in the dilepton pair center of mass. The two hadrons momenta $\vec{p}'_A$ and $\vec{p}'_B$ span the orange plane and the z -axis is defined as the bisector of $\vec{p}'_A$ and $-\vec{p}'_B$. The two leptons are emitted back to back in the black plane. In addition, in symmetric proton-proton collisions, proton A is defined such as it is the proton flying in the same direction than the dilepton longitudinal boost in the laboratory frame. With such a choice, in the case of a $q\bar{q}$ annihilation, θ_{CS} provides an approximation of θ defined as the angle between the negative lepton and the quark.

7.2.3 The Drell-Yan process and spin 1 resonances

The Drell-Yan angular differential cross section is very well known and is, at leading order, of the form¹:

$$\frac{d\sigma_{\gamma/Z}}{d\Omega} \propto c_1(1 + \cos^2 \theta) + c_2 \cos \theta \quad (7.7)$$

where c_1 and c_2 are two parameters that contain the different couplings and depend on the dilepton invariant mass.

This formula holds for any process of the type $q\bar{q} \rightarrow S^1 \rightarrow l\bar{l}$ where S^1 is a spin 1 particle. This means that both c_1 and c_2 must be measured to be able to use the angular distribution to distinguish between two spin 1 models. The term c_2 , which is responsible for the forward-backward asymmetry, is impossible to measure without the knowledge of the quark and lepton directions. This discussion is illustrated in figure 7.4 which shows the distribution, at the generated level and without any acceptance cut, of $\cos \theta_{CS}$ (left) and $|\cos \theta_{CS}|$ (right) for two Z' models (Z'_ψ and Z'_{SSM}) generated at a mass of 3 TeV and Drell-Yan events with a mass between 2.7 and 3.3 TeV.

As discussed in section 1.2.2, the Drell-Yan process shows a strong positive forward-backward asymmetry for dilepton masses above 120 GeV/ c^2 . This allows one to differentiate it from the two signal models. All the discrimination power of $\cos \theta_{CS}$ vanishes once the absolute value is taken, i.e. if one does not distinguish the particles from the antiparticles. On the right plot, a fit to the function $(1 + \cos^2 \theta_{CS})$ is also performed. The good performance of the fit indicates that $\cos \theta_{CS}$ provides indeed a good approximation of $\cos \theta$.

¹cf. Appendix A.

7. STUDY OF THE CHARACTERISATION OF A POTENTIAL NEW DIELECTRON RESONANCE SIGNAL THAT COULD BE DISCOVERED IN THE EARLY DATA AT $\sqrt{s} = 13$ AND 14 TEV

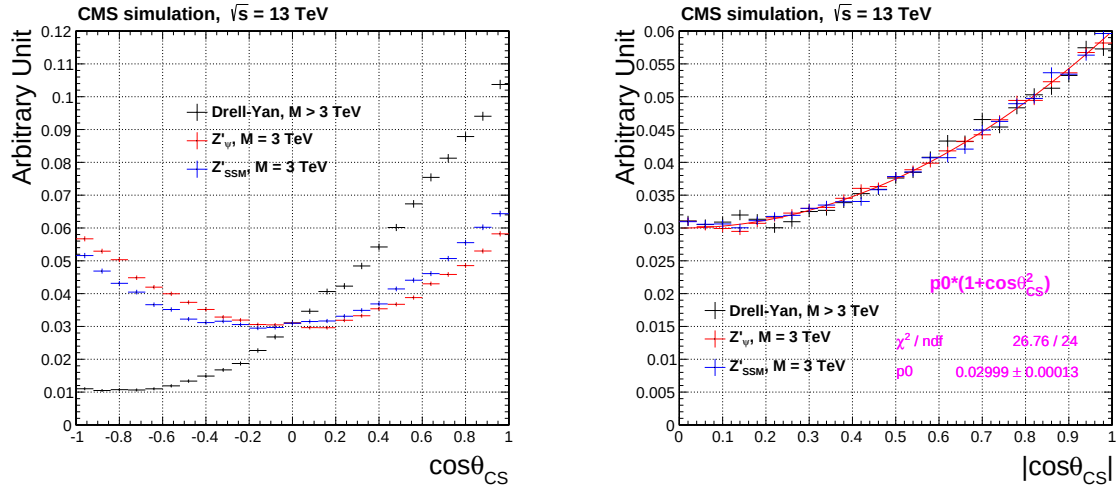


Figure 7.4: The distribution of $\cos\theta_{CS}$ (left) and $|\cos\theta_{CS}|$ (right) at the generated level for two different Z' models with a mass of 3 TeV and for Drell-Yan events with an invariant mass above 3 TeV. The identical $|\cos\theta_{CS}|$ distributions for the three models indicate that the distinction between several models where the exchanged particle carries a spin 1 can only be made if one knows the quark and lepton directions unambiguously. On the right plot, the good agreement of the fit of the distribution to the function $1 + \cos^2\theta_{CS}$ indicates that θ_{CS} approximates well θ .

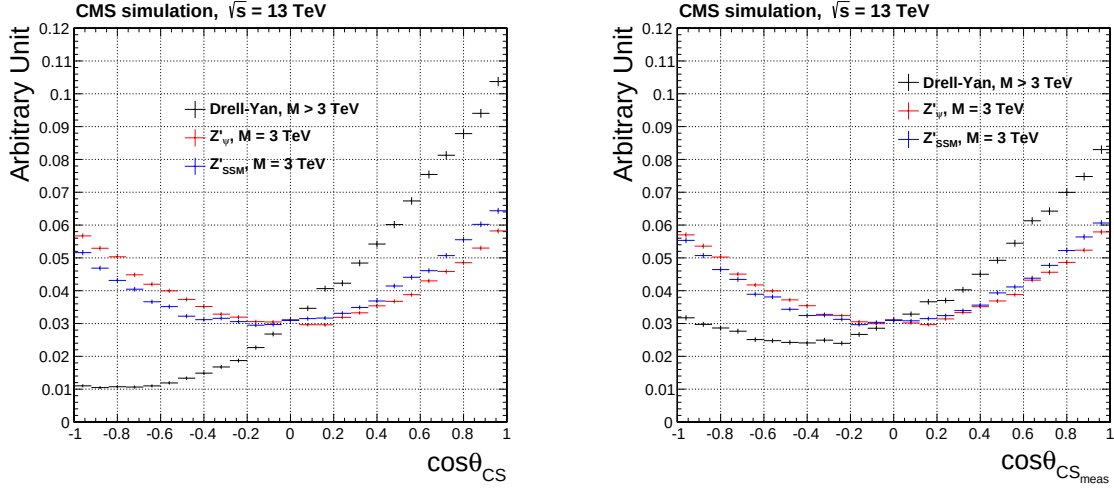


Figure 7.5: The distribution of $\cos\theta_{CS}$ (left) and $\cos\theta_{CS_{meas}}$ (right) at the generated level for two different Z' models with a mass of 3 TeV and for Drell-Yan events with an invariant mass above 3 TeV. The forward-backward asymmetry is diluted in the measured distribution because of the uncertainty on the quark direction.

7.2.4 Identification of the quark direction

In the previous section, the results shown used the $\cos\theta_{CS}$ variable where the quark direction is known. The experimentally accessible variable however is $\cos\theta_{CS_{meas}}$. The two variables, at the generated level and without any acceptance cut, are compared in figure 7.5. For events where the quark is less energetic than the antiquark, the assumption on the quark direction is uncorrect and $\cos\theta_{CS_{meas}} = -\cos\theta_{CS}$. The resulting effect is a dilution of the asymmetry in the $\cos\theta_{CS_{meas}}$ distribution. The probability of misidentifying the quark direction in Drell-Yan events depends on the magnitude of the longitudinal boost of the dilepton pair. For a boost tending to 0, this probability will tend to 50%, while a very large boost implies that the responsible parton carries a large fraction of the proton energy which increases its probability to come from the valence. It is therefore interesting to study this dependency. Instead of studying directly the boost, the rapidity of the pair, y_{ll} , is used. The performances of the two variables are exactly identical.

Figure 7.6 presents the probability of correctly identifying the quark direction as a function of the fraction of selected events when a condition $|y_{ll}| > y_{cut}$ is applied. The Drell-Yan simulated samples mentioned in table 3.1 and produced at $\sqrt{s} = 8$ TeV were used for this study. Some cut values y_{cut} are indicated in red. Several mass ranges were considered, from the Z peak region up to 3 TeV. Overall, the probability increases with the mass, up to 2 TeV, for any fraction of selected events. High masses increase the average proton energy fraction carried by each of the parton and increases, for the quarks, the valence contribution. A slight drop is however observed for events between 2 and 3 TeV. Up to mass of 1 TeV, the cut values to reach a given efficiency are slowly evolving. Beyond 1 TeV, they are strongly moving.

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In summary, the probability of misidentifying the quark direction shows a clear dependence on the dilepton rapidity. The use of this variable brings therefore an additional information that can improve the precision of a study of the lepton angular distribution.

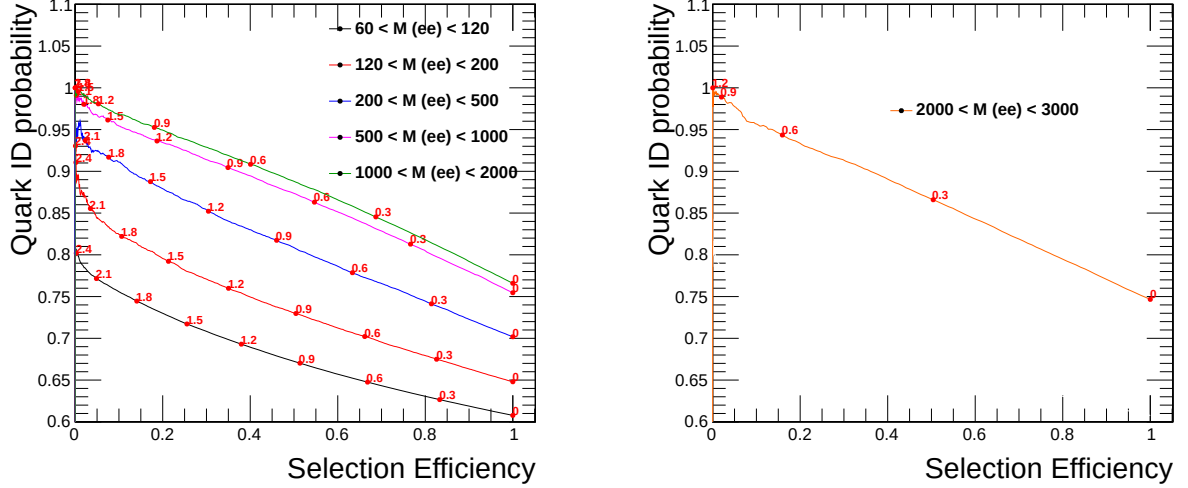


Figure 7.6: Probability of correctly identifying the quark direction as a function of the fraction of selected events when a condition $|y_{l\bar{l}}| > y_{cut}$ is applied for different mass ranges up to 2 TeV (left) and for masses from 2 to 3 TeV (right) for Drell-Yan simulated events at $\sqrt{s} = 8$ TeV. Several cut values y_{cut} are indicated in red.

7.3 Results

This section presents the measurement of three characteristics of a new resonance: its spin, its forward-backward asymmetry and its production modes. In the following, three benchmark models are considered (the Z'_ψ and the G_{RS} for the spin determination, the Z'_ψ and Z'_{SSM} for the forward-backward asymmetry and the G_{RS} for the production mode fractions measurement) and their properties are estimated focusing on techniques as model-independent as possible.

7.3.1 Spin measurement

The final state dilepton angular distribution for the $parton\ parton \rightarrow l\bar{l}$ process depends on the spin of the exchanged particle and on the production mode. A spin 0 particle can be produced through gg or $q\bar{q}$ fusion and will lead to a flat $\cos\theta$ distribution. For a spin 1 or 2 particle, the $\cos\theta$ is distributed according to the sum of an even polynomial, whose coefficients are univoquely determined for a given production mode ($q\bar{q}$ for spin 1¹, gg or

¹The Landau-Yang theorem states that a spin 1 resonance cannot be produced from two identical massless spin 1 particles [135, 136, 137].

resonance spin and production mode	$d\sigma/d \cos\theta $
Spin 0 (gg or $q\bar{q}$ fusion)	$\propto 1$
Spin 1 ($q\bar{q}$ fusion)	$\propto (1 + \cos^2\theta)$
Spin 2 (gg fusion)	$\propto (1 - \cos^4\theta)$
Spin 2 ($q\bar{q}$ fusion)	$\propto (1 - 3\cos^2\theta + 4\cos^4\theta)$

Table 7.3: The $|\cos\theta|$ dependency of the differential cross section ($d\sigma/d|\cos\theta|$) for the process $parton\ parton \rightarrow l\bar{l}$ for different resonance spins and production modes. In the case of a spin 1 or 2 resonance, a non zero forward-backward asymmetry can arise through terms proportional to odd powers of $\cos\theta$.

$q\bar{q}$ for a spin 0 or 2) and an odd polynomial, that leads to a forward-backward asymmetry and which is model dependent. In order to be as model independent as possible and to get rid of the asymmetry, only the absolute value of $\cos\theta$ will be used in the spin measurement.

The $|\cos\theta|$ dependency [138, 139] is summarized in table 7.3 for the different spin scenarios and production modes. The production mode has several impact on the experimental variable $\cos\theta_{CS_{meas}}$.

As already mentioned, the two production modes lead to different $\cos\theta$ shapes for a spin 2, but they also induce different event kinematics at a proton-proton collider. At moderate x values, the contribution from the valence quarks is sizeable (*cf.* section 3.1.1). Those quarks carry in average more momentum than the antiquarks and the $q\bar{q}$ system is boosted along the z-axis in the quark direction. One therefore expects the two leptons produced from a gg fusion to be more central than those coming from a $q\bar{q}$ fusion. A difference is also expected between up and down quarks since the valence contribution is twice larger in the first case compared to the second.

It has been checked that the impact of the dilepton pair longitudinal boost on the acceptance times efficiency is small (below 1%). However, the dilepton rapidity distributions show sizeable differences. The use of this variable can help to enhance the separation between different spin scenarios but introduces some model dependency.

In this section the expected separation between a signal hypothesis and various alternative hypotheses is measured. Two cases are considered: a Z'_ψ and a G_{RS} . Their $\cos\theta_{CS}$ distributions are shown in figure 7.7, separately for the gg and $q\bar{q}$ production modes in the case of the graviton.

The Z'_ψ and the RS graviton have the following properties in common:

- Similar cross sections
- Same coupling to up and down quarks
- $A_{FB} = 0$ (if the interference with the Drell-Yan process is neglected)

The alternative hypotheses are listed in table 7.4. For the spin 0 and 2, various mixtures of gg and $q\bar{q}$ fractions are considered. The distributions for those hypotheses are built from the Randall Sundrum and Z'_ψ distributions reweighted according to the desired gg and $q\bar{q}$ fractions and $\cos\theta_{CS}$ distribution.

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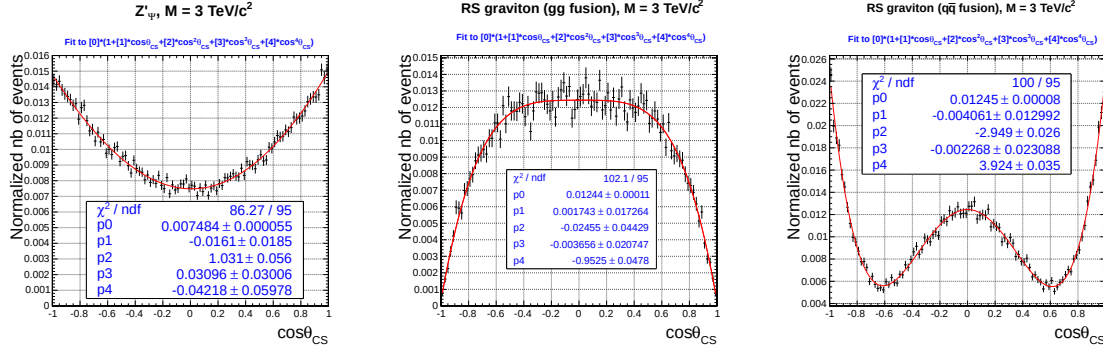


Figure 7.7: Normalized $\cos\theta_{CS}$ distributions for the two signals considered, namely the Z'_ψ (left) and the G_{RS} , according to the production mode (gg (center) or $q\bar{q}$ (right)) for the graviton. A fit to a fourth order polynomial is also performed, giving compatible results with the theoretical predictions for $\cos\theta$ (table 7.3).

Spin (production mode)	Construction
Spin 0 (100% $q\bar{q}$, 50% gg -50% $q\bar{q}$, 100% gg)	Reweighting the $ \cos\theta_{CS} $ and gg - $q\bar{q}$ distributions of RS graviton.
Spin 0 (50% gg -50% $q\bar{q}$)	
Spin 0 (100% gg)	
Spin 1 (100% $q\bar{q}$)	Using the Z'_ψ distribution.
Spin 2 (100% $q\bar{q}$)	Reweighting the RS graviton gg - $q\bar{q}$ fractions.
Spin 2 (50% gg -50% $q\bar{q}$)	
Spin 2 (100% gg)	

Table 7.4: Alternative hypotheses for the spin measurement.

7.3.1.1 Procedure

The method used here is the CL_s method [140] that was already used for measuring the property of the new 125 GeV particle recently discovered [141]. It consists of performing a test between two hypotheses H_1 (signal) and H_0 (alternative hypothesis). For a given set of n observed events each one characterized by k variables and denoted $\{\vec{x}_1, \vec{x}_2, \dots, \vec{x}_n\}$ where $\vec{x}_i$ has k components and has a probability density function (pdf) $p_0(\vec{x})$ ($p_1(\vec{x})$) under the hypothesis H_0 (H_1), one defines the log likelihood ratio, Q , as:

$$Q = \ln \frac{\mathcal{L}_0}{\mathcal{L}_1} = \ln \frac{\prod_{i=1}^n p_0(\vec{x}_i)}{\prod_{i=1}^n p_1(\vec{x}_i)} \quad (7.8)$$

It can be shown that this test statistics is unbiased and presents the best discrimination power between the two hypotheses. Let us denote q_0 (q_1) the pdf for Q obtained if H_0 (H_1) is true. In practice this pdf is obtained by generating a large number of pseudo-experiments under H_1 or H_0 and calculating Q for each of them. If $Q_{obs} < 0$, the hypothesis

H_0 is disfavoured with respect to H_1 in the data with a CL_s value given by:

$$CL_s = \frac{\int_{-\infty}^{Q_{obs}} q_0 dQ}{\int_{-\infty}^{Q_{obs}} q_1 dQ} \quad (7.9)$$

The expected CL_s value in case H_1 is true can be obtained by replacing Q_{obs} by the median value of the q_1 distribution. It is converted into a number of separation in standard deviations (σ) given by $\sqrt{2}\text{Erf}^{-1}(1 - CL_s)$ (where Erf^{-1} is the inverse function associated to the error function).

In this analysis, the pdf $p_0(\vec{x})$ and $p_1(\vec{x})$ are binned and obtained from the Monte Carlo simulation. In order to avoid a possible bias due to the limited statistics, separate events are used to build the pdf used for the pseudo-experiments generation and those used in the likelihood calculation.

7.3.1.2 Results

The expected separation between different spin hypotheses and production modes are presented for likelihood based on a 1 dimensional (1D) or 2 dimensional (2D) pdf in figures 7.8 and 7.9 and tables 7.5 and 7.6 for the three luminosity scenarios mentioned above. The 1D pdf only uses the $|\cos\theta_{CS_{meas}}|$ distribution while the dielectron pair rapidity y_{ee} is added in the 2D pdf. Both pdfs contain the same number of bins (100 bins in $|\cos\theta_{CS_{meas}}|$ for the 1D pdf, 20 bins in $|\cos\theta_{CS_{meas}}|$ times 5 bins in $|y_{ee}|$ for the 2D pdf).

For a Z'_ψ or a G_{RS} at 3 TeV, the alternative hypotheses are typically excluded by 1σ to 3σ for 100 fb^{-1} . A 5σ separation can be achieved with the full dataset of the LHC phase 2. The only exception is the spin 2 ($q\bar{q}$ - gg) hypothesis which is close to the G_{RS} signal for which the gg and $q\bar{q}$ production fractions are 0.38 and 0.62 respectively.

For the 4 TeV signals the full dataset of the LHC phase 2 is required to exclude all the alternative hypotheses at the 2σ level or more. As discussed below, the gg and $q\bar{q}$ production fractions for the G_{RS} are mass dependent. The $q\bar{q}$ fraction rises to 0.75 for 4 TeV graviton at $\sqrt{s} = 13 \text{ TeV}$. This explains why the spin 2 ($q\bar{q}$ - gg) hypothesis is less disfavoured compared to the spin 2 ($q\bar{q}$) hypothesis for $M_{G_{RS}} = 3 \text{ TeV}$ whereas both options are equally disfavoured at $M_{G_{RS}} = 4 \text{ TeV}$.

In most of the cases, the likelihood based on the 2D pdf significantly improves the discrimination compared to the 1D pdf for hypotheses or signals where the two production modes are involved.

An example of a pseudo-experiment generated according to the G_{RS} ($M=3 \text{ TeV}$) hypothesis at 3000 fb^{-1} is shown in figure 7.10 and figure 7.11 for the 1D and 2D likelihood analyses respectively. The G_{RS} and Z'_ψ pdfs multiplied by the number of expected events are also drawn. The discrimination offered by the use of the dilepton pair rapidity is clearly visible in figure 7.11.

The 1D and 2D likelihood ratios for a G_{RS} tested against a spin 1 ($q\bar{q}$) are presented in figure 7.12 and figure 7.13 for the two mass points and the 300 and 3000 fb^{-1} luminosity scenarios.

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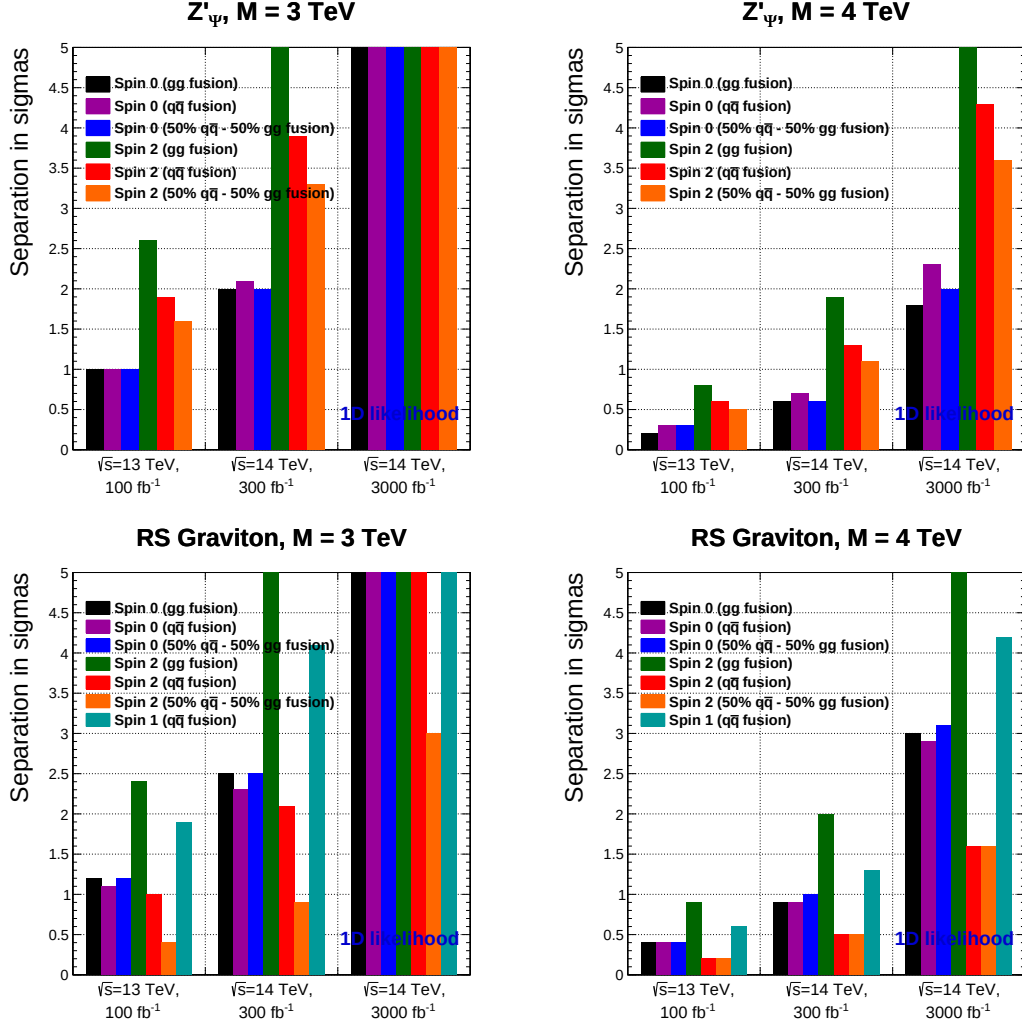


Figure 7.8: Expected separation (in σ) between different spin hypotheses for three luminosity scenarios. Two signal hypotheses (H_1) are considered: a Z'_ψ (top) and a G_{RS} (bottom) with a mass of 3 (left) or 4 (right) TeV. They are compared to various alternative hypotheses (H_0), shown by the various curves. The results are based on a likelihood using a single variable: $|\cos \theta_{CS}|$.

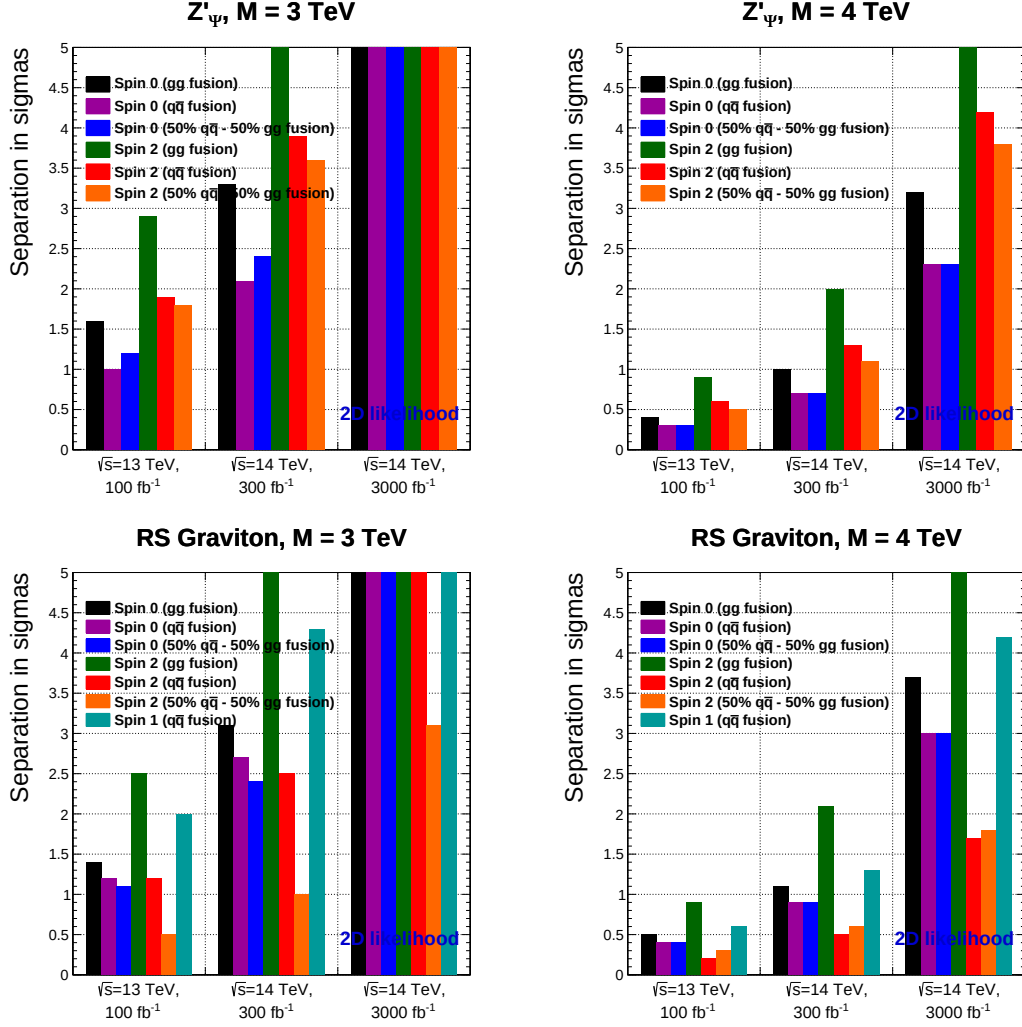


Figure 7.9: Expected separation (in σ) between different spin hypotheses for three luminosity scenarios. Two signal hypotheses (H_1) are considered: a Z'_ψ (top) and a G_{RS} (bottom) with a mass of 3 (left) or 4 (right) TeV. They are compared to various alternative hypotheses (H_0), shown by the various curves. The results are based on a likelihood using two variable: $|\cos \theta_{CS}|$ and $|y_{ee}|$.

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H_1	H_0	100 fb ⁻¹ (13 TeV)	300 fb ⁻¹ (14 TeV)	3000 fb ⁻¹ (14 TeV)
Z'_ψ M=3 TeV	Spin 0 (gg)	1.0	2.0	> 5
Z'_ψ M=3 TeV	Spin 0 ($q\bar{q}$)	1.0	2.1	> 5
Z'_ψ M=3 TeV	Spin 0 ($q\bar{q}$ - gg)	1.0	2	> 5
Z'_ψ M=3 TeV	Spin 2 (gg)	2.6	> 5	> 5
Z'_ψ M=3 TeV	Spin 2 ($q\bar{q}$)	1.9	3.9	> 5
Z'_ψ M=3 TeV	Spin 2 ($q\bar{q}$ - gg)	1.6	3.3	> 5
Z'_ψ M=4 TeV	Spin 0 (gg)	0.2	0.6	1.8
Z'_ψ M=4 TeV	Spin 0 ($q\bar{q}$)	0.3	0.7	2.3
Z'_ψ M=4 TeV	Spin 0 ($q\bar{q}$ - gg)	0.3	0.6	2.0
Z'_ψ M=4 TeV	Spin 2 (gg)	0.8	1.9	> 5
Z'_ψ M=4 TeV	Spin 2 ($q\bar{q}$)	0.6	1.3	4.3
Z'_ψ M=4 TeV	Spin 2 ($q\bar{q}$ - gg)	0.5	1.1	3.6
G_{RS} M=3 TeV	Spin 0 (gg)	1.2	2.5	> 5
G_{RS} M=3 TeV	Spin 0 ($q\bar{q}$)	1.1	2.3	> 5
G_{RS} M=3 TeV	Spin 0 ($q\bar{q}$ - gg)	1.2	2.5	> 5
G_{RS} M=3 TeV	Spin 2 (gg)	2.4	> 5	> 5
G_{RS} M=3 TeV	Spin 2 ($q\bar{q}$)	1.0	2.1	> 5
G_{RS} M=3 TeV	Spin 2 ($q\bar{q}$ - gg)	0.4	0.9	3.0
G_{RS} M=3 TeV	Spin 1 ($q\bar{q}$)	1.9	4.1	> 5
G_{RS} M=4 TeV	Spin 0 (gg)	0.4	0.9	3.0
G_{RS} M=4 TeV	Spin 0 ($q\bar{q}$)	0.4	0.9	2.9
G_{RS} M=4 TeV	Spin 0 ($q\bar{q}$ - gg)	0.4	1	3.1
G_{RS} M=4 TeV	Spin 2 (gg)	0.9	2.0	> 5
G_{RS} M=4 TeV	Spin 2 ($q\bar{q}$)	0.2	0.5	1.6
G_{RS} M=4 TeV	Spin 2 ($q\bar{q}$ - gg)	0.2	0.5	1.6
G_{RS} M=4 TeV	Spin 1 ($q\bar{q}$)	0.6	1.3	4.2

Table 7.5: Expected separation (in σ) between different spin hypotheses for the three luminosity scenarios. Two signal hypotheses (H_1) are considered: a Z'_ψ and a G_{RS} with a mass of 3 or 4 TeV. They are compared to various alternative hypotheses (H_0). The results are based on a likelihood using a single variable: $|\cos \theta_{CS}|$.

H_1	H_0	100 fb ⁻¹ (13 TeV)	300 fb ⁻¹ (14 TeV)	3000 fb ⁻¹ (14 TeV)
Z'_ψ M=3 TeV	Spin 0 (gg)	1.6	3.3	> 5
Z'_ψ M=3 TeV	Spin 0 ($q\bar{q}$)	1.0	2.1	> 5
Z'_ψ M=3 TeV	Spin 0 ($q\bar{q}$ - gg)	1.2	2.4	> 5
Z'_ψ M=3 TeV	Spin 2 (gg)	2.9	> 5	> 5
Z'_ψ M=3 TeV	Spin 2 ($q\bar{q}$)	1.9	3.9	> 5
Z'_ψ M=3 TeV	Spin 2 ($q\bar{q}$ - gg)	1.8	3.6	> 5
Z'_ψ M=4 TeV	Spin 0 (gg)	0.4	1	3.2
Z'_ψ M=4 TeV	Spin 0 ($q\bar{q}$)	0.3	0.7	2.3
Z'_ψ M=4 TeV	Spin 0 ($q\bar{q}$ - gg)	0.3	0.7	2.3
Z'_ψ M=4 TeV	Spin 2 (gg)	0.9	2	> 5
Z'_ψ M=4 TeV	Spin 2 ($q\bar{q}$)	0.6	1.3	4.2
Z'_ψ M=4 TeV	Spin 2 ($q\bar{q}$ - gg)	0.5	1.1	3.8
G_{RS} M=3 TeV	Spin 0 (gg)	1.4	3.1	> 5
G_{RS} M=3 TeV	Spin 0 ($q\bar{q}$)	1.2	2.7	> 5
G_{RS} M=3 TeV	Spin 0 ($q\bar{q}$ - gg)	1.1	2.4	> 5
G_{RS} M=3 TeV	Spin 2 (gg)	2.5	> 5	> 5
G_{RS} M=3 TeV	Spin 2 ($q\bar{q}$)	1.2	2.5	> 5
G_{RS} M=3 TeV	Spin 2 ($q\bar{q}$ - gg)	0.5	1	3.1
G_{RS} M=3 TeV	Spin 1 ($q\bar{q}$)	2.0	4.3	> 5
G_{RS} M=4 TeV	Spin 0 (gg)	0.5	1.1	3.7
G_{RS} M=4 TeV	Spin 0 ($q\bar{q}$)	0.4	0.9	3
G_{RS} M=4 TeV	Spin 0 ($q\bar{q}$ - gg)	0.4	0.9	3.1
G_{RS} M=4 TeV	Spin 2 (gg)	0.9	2.1	> 5
G_{RS} M=4 TeV	Spin 2 ($q\bar{q}$)	0.2	0.5	1.7
G_{RS} M=4 TeV	Spin 2 ($q\bar{q}$ - gg)	0.3	0.6	1.8
G_{RS} M=4 TeV	Spin 1 ($q\bar{q}$)	0.6	1.3	4.2

Table 7.6: Expected separation (in σ) between different spin hypotheses for three luminosity scenarios. Two signal hypotheses (H_1) are considered: a Z'_ψ and a G_{RS} with a mass of 3 or 4 TeV. They are compared to various alternative hypotheses (H_0). The results are based on a likelihood using two variables: $|\cos \theta_{CS}|$ and $|y_{ee}|$.

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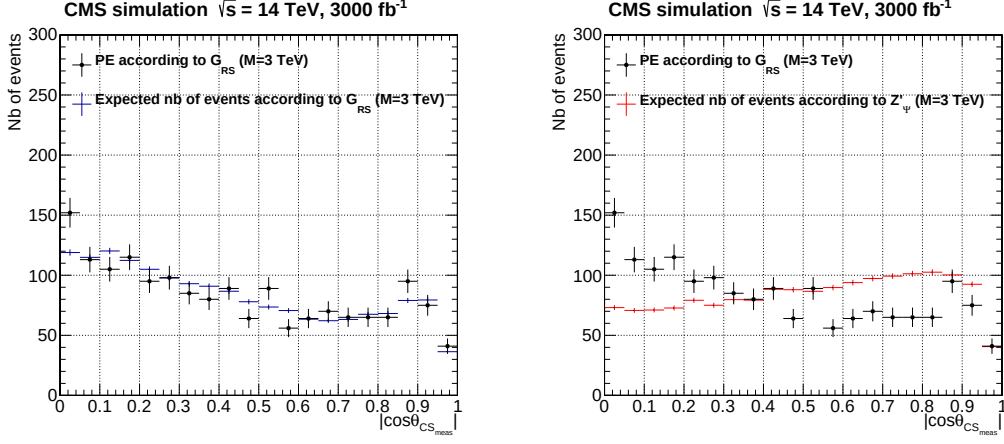


Figure 7.10: Example of a pseudo-experiment generated according to the G_{RS} (M=3 TeV) $\cos\theta_{CS}$ distribution at 3000 fb^{-1} . The pseudo data are shown on top of the G_{RS} (left) and Z'_ψ (right) pdfs multiplied by the number of expected events.

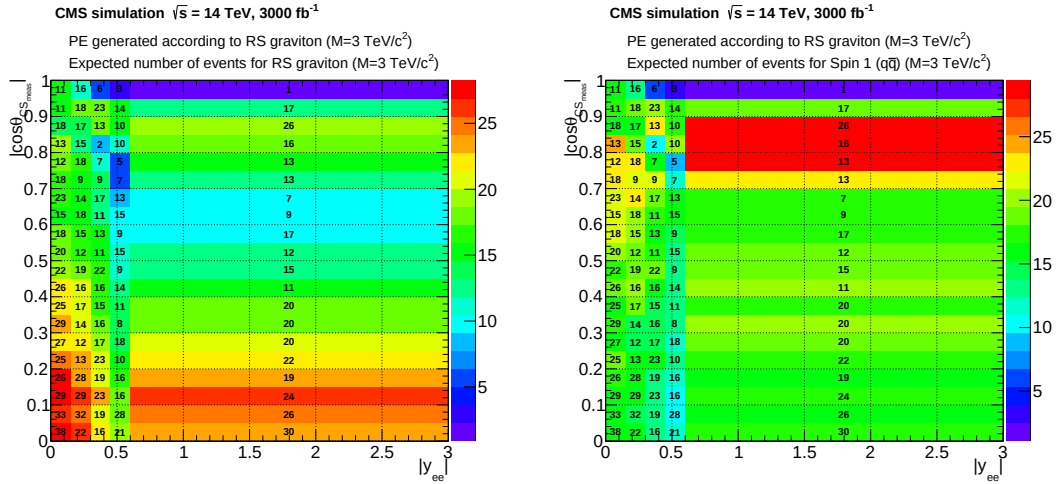


Figure 7.11: Example of a pseudo-experiment generated according to the G_{RS} (M=3 TeV) ($|\cos\theta_{CS}|, |y_{ee}|$) distribution at 3000 fb^{-1} . The pseudo data (numbers in the bin) are shown on top of the G_{RS} (left) and Z'_ψ (right) pdfs multiplied by the number of expected events. In the $|y_{ee}|$ coordinate, all the bins have the same size except the last bin which goes from 0.6 to 3.

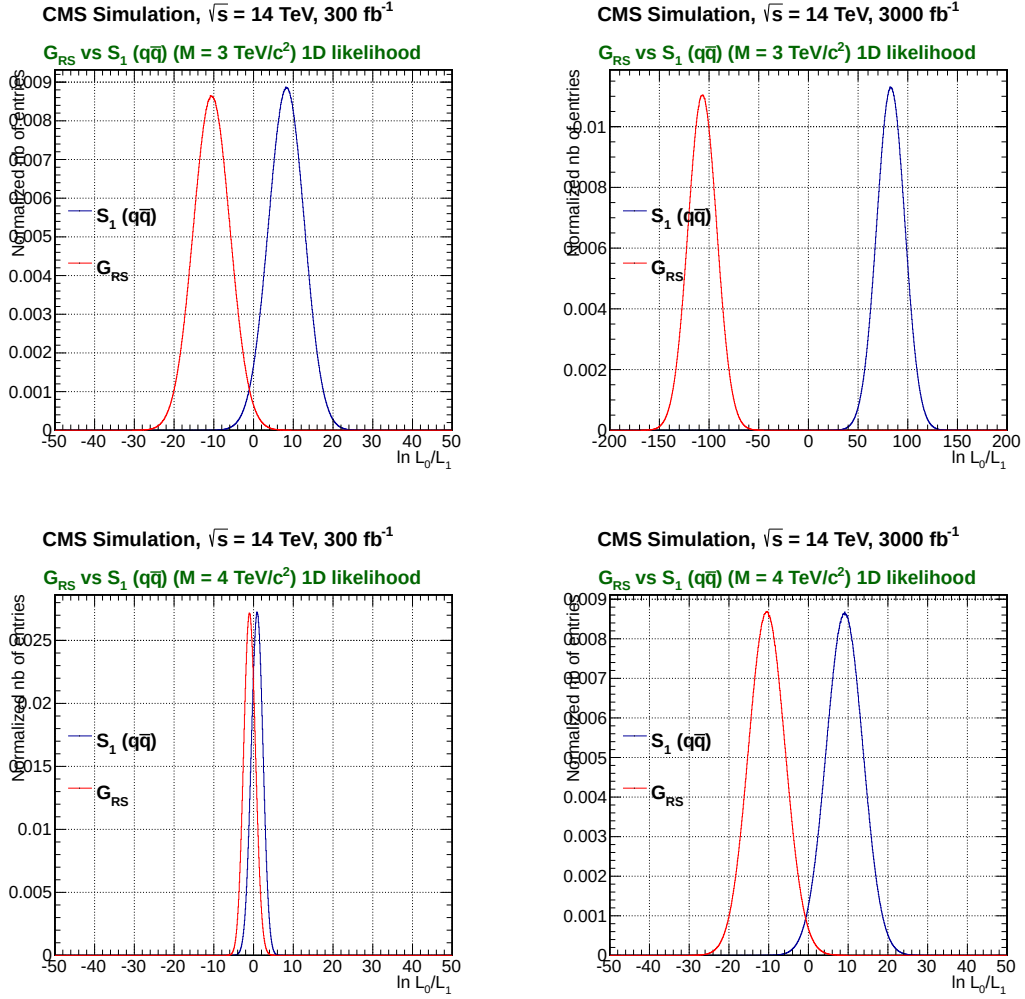


Figure 7.12: Likelihood ratios obtained from the $|\cos \theta_{CS}|$ distribution alone for a G_{RS} tested against a spin 1 ($q\bar{q}$). Two luminosity scenarios and two mass points (3 TeV, top, and 4 TeV, bottom) are presented.

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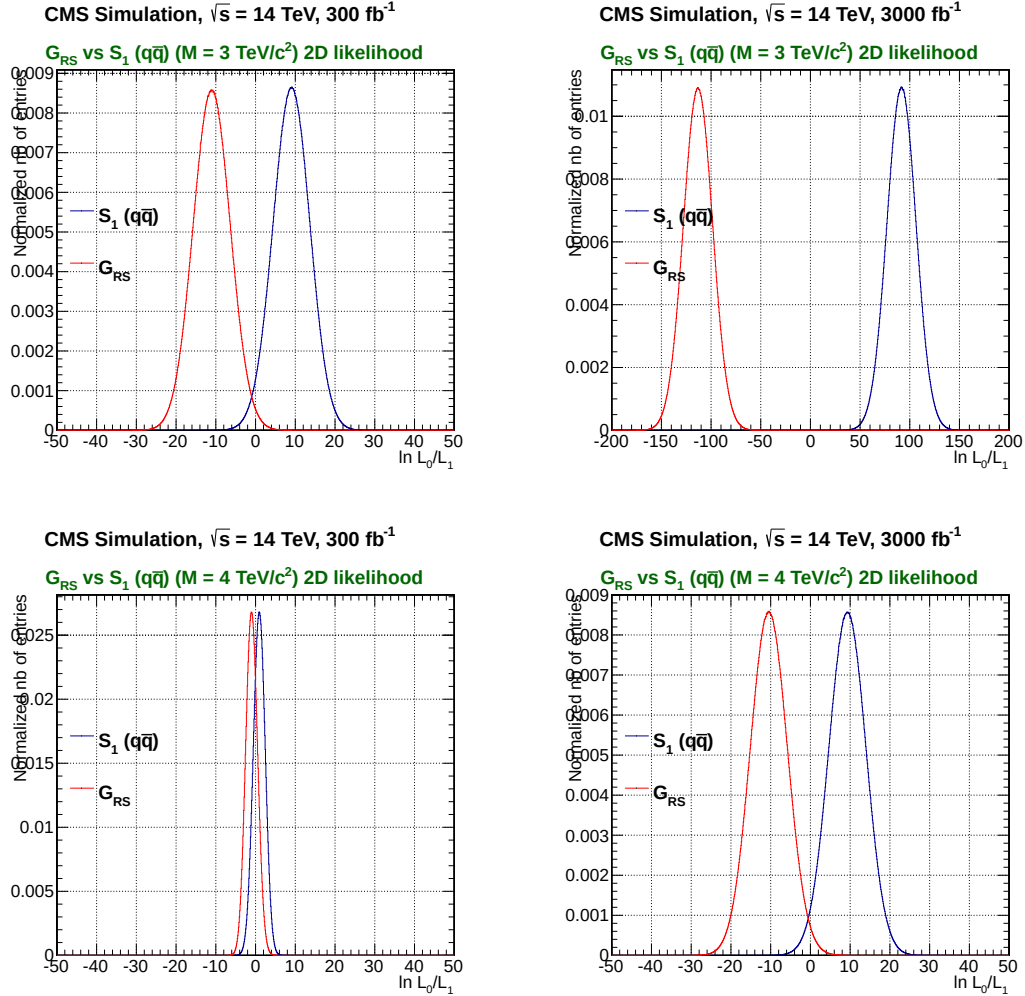


Figure 7.13: Likelihood ratios obtained from the $(|\cos\theta_{CS}|, |y_{ee}|)$ distribution for a G_{RS} tested against a spin 1 ($q\bar{q}$). Two luminosity scenarios and two mass points (3 TeV, top, and 4 TeV, bottom) are presented.

7.3.2 A_{FB} measurement

In case a new resonance is found, another interesting property to study is the forward-backward asymmetry. As discussed in section 1.4.2.2, this variable is extremely model dependent and can therefore provide valuable information to understand the newly discovered particle. Unlike the spin measurement, the A_{FB} measurement requires the correct identification of the quark and the negative lepton. Only events with two opposite charge electrons are therefore selected for this measurement. As mentioned in section 7.2, the quark direction identification is subject to measurement errors: the wrong direction is chosen when the antiquark is more energetic than the quark. The use of the dielectron rapidity reduces this uncertainty and is therefore expected to improve the precision on the result.

7.3.2.1 Procedure

The study presented here is limited to the measurement of the expected forward-backward asymmetry and its uncertainty for several spin 1 resonances. The measurement is achieved through a likelihood minimization method. For a spin 1 resonance, since the differential cross section in equation 7.7 must be defined positive, A_{FB} must lie in the range $[-0.75, +0.75]$. Two samples corresponding to two extreme cases ($A_{FB} = \pm 0.75$) are therefore built by reweighting the $\cos \theta_{CS}$ distribution of the Z'_ψ signal. The Z'_ψ has the particular feature of equal coupling to up and down quarks. The $\cos \theta_{CS}$ distribution at the generated level and the $\cos \theta_{CS_{meas}}$ after the complete event selection are compared in figure 7.14 for the two extreme cases. The drop of events at high $\cos \theta_{CS_{meas}}$ is due to a kinematical effect: the negative electron in events with $\cos \theta_{CS}$ values close to 1 are very forward and sometimes emitted at $|\eta|$ outside the tracker acceptance. The 2-dimensional distribution of $(\cos \theta_{CS_{meas}}, |y_{ee}|)$ is shown in figure fig:15 and illustrates the fact that the identification of the quark direction is more efficient for high rapidity values.

Let $p_{A_{FB}=\pm 0.75}$ be the corresponding pdfs for the experimental variables $\vec{x}$ (here $\cos \theta_{CS_{meas}}$ or $(\cos \theta_{CS_{meas}}, |y_{ee}|)$). Any spin 1 distribution is then a linear combination of the two extreme cases:

$$p_\kappa = \kappa \times p_{A_{FB}=+0.75} + (1 - \kappa) \times p_{A_{FB}=-0.75} \quad (7.10)$$

characterized by a number κ related to the forward-backward asymmetry¹:

$$A_{FB} = \frac{3}{4}(2\kappa - 1) \quad (7.11)$$

Pseudo experiments are then generated according to the signal under study and the κ value that maximizes the likelihood $\prod_{i=1}^n p_\kappa(\vec{x}_i)$ is estimated. The median and the $\pm 1\sigma$ bands are then obtained by repeating this procedure 1000 times.

¹ $\kappa \times p_{A_{FB}=+0.75} + (1 - \kappa)p_{A_{FB}=-0.75}$
 $\propto \kappa(1 + \cos^2 \theta + 2 \cos \theta) + (1 - \kappa)(1 + \cos^2 \theta - 2 \cos \theta)$
 $\propto 1 + \cos^2 \theta + (4\kappa - 2) \cos \theta$
 $\propto 1 + \cos^2 \theta + \frac{8}{3}A_{FB} \cos \theta$

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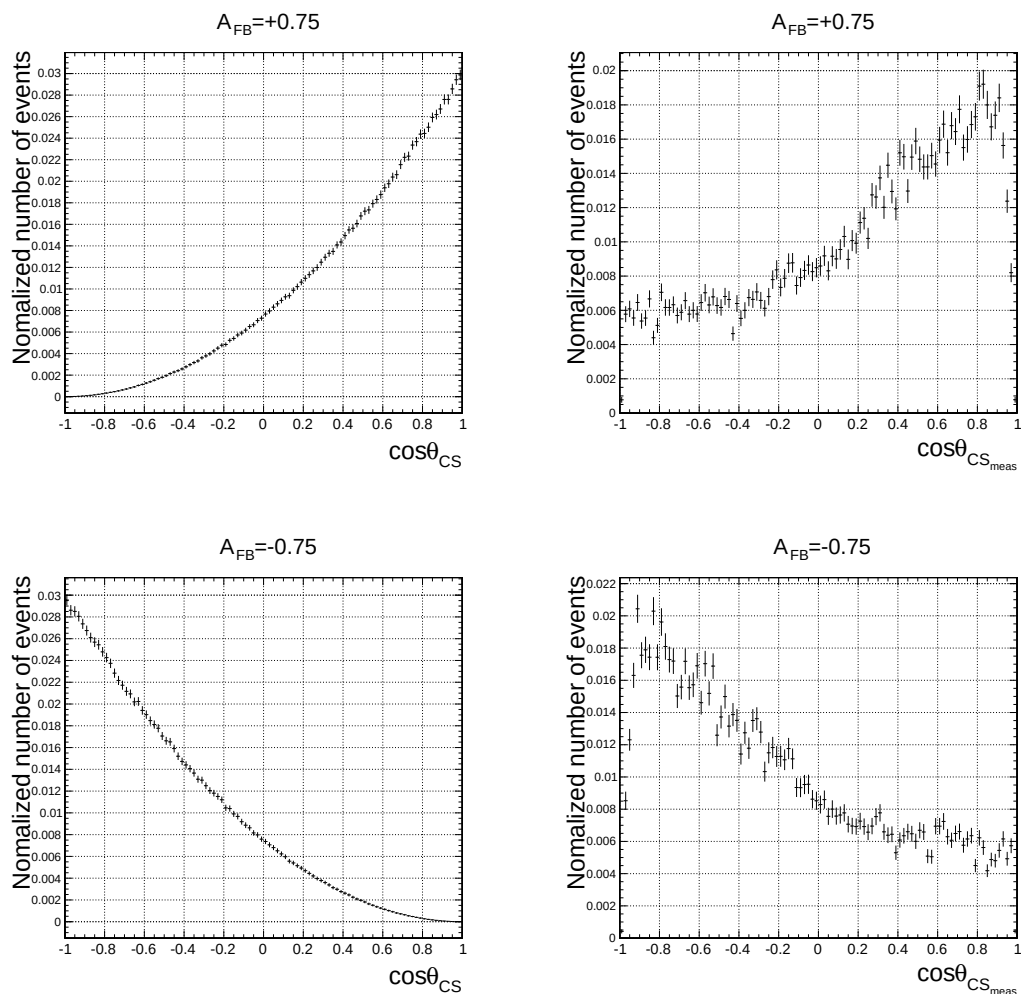


Figure 7.14: Construction of the $A_{FB} = +0.75$ (top) or -0.75 (bottom) distributions. Events from Z'_ψ are reweighted such that their $\cos\theta_{CS}$ distribution (left) is proportional to $1 + \cos^2\theta_{CS} \pm 2\cos\theta_{CS}$. The distribution of the measured variable for events passing the selection, $\cos\theta_{CS_{meas}}$, can then be estimated (right).

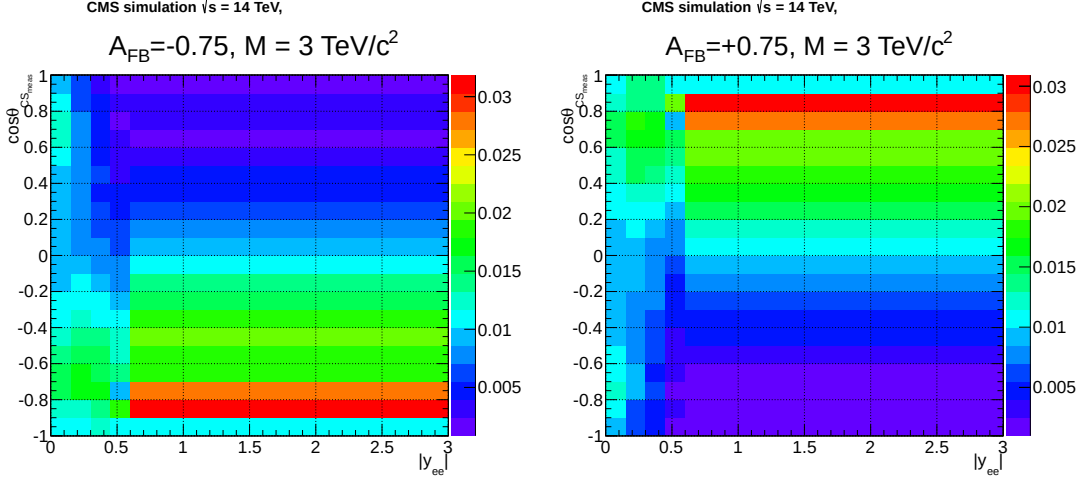


Figure 7.15: $(\cos \theta_{CS_{meas}}, |y_{ee}|)$ distributions for the two extreme cases $A_{FB} = \pm 0.75$. In the $|y_{ee}|$ coordinate, all the bins have the same size except the last bin which goes from 0.6 to 3. The samples are built by reweighting Z'_ψ events. The $\cos \theta_{CS_{meas}}$ distribution is rapidity dependent: the quark direction is in average better identified for large $|y_{ee}|$.

7.3.2.2 Results

The results for a Z'_{SSM} and a Z'_ψ are shown in figures 7.16 and 7.17 and in tables 7.7 and 7.8. The theoretical values of the forward-backward asymmetry for those signals and for the Drell-Yan process are listed in table 7.9.

The small discrepancies between the median values and the theoretical expectations is explained by the finite statistics in the simulated samples.

In this study, the quantity of interest is the size of the 1σ band, which represents the expected statistical precision that can be reached on the forward-backward asymmetry measurement. It follows a $1/\sqrt{N}$ law, where N is the number of signal events. For the full luminosity delivered during the LHC phase 2, a precision of ± 0.05 for 3 TeV signals and ± 0.1 - 0.15 for the 4 TeV signals is reached.

The two signals studied here have a small or null asymmetry. However, as shown in figure 1.5 of section 1.4.2.2, signals with any asymmetry exist. An example is the case of the Z'_I which only couples to left-handed electrons and right-handed down quarks and where the asymmetry is maximally negative. For masses above 200 GeV, the Drell-Yan process is for its part characterized by a strong positive asymmetry. For comparison purposes, in tables 7.7 and 7.8, results are therefore also presented for 1056 Drell-Yan events, in which a strong forward-backward asymmetry is present. This number of events corresponds to the 3 TeV Z'_ψ yield for 3000 fb^{-1} . This shows that the error on the measurement is not sensitive to its central value.

As expected, the 2D likelihood slightly improves the precision of the measurement. This 2D method is however expected to show some model dependency: for the likelihood calculation the $|y_{ee}|$ distribution from Z'_ψ events was used, which slightly differs from the Z'_{SSM} or Drell-Yan one. Figure 7.18 shows the likelihood based on the 1D or 2D pdf

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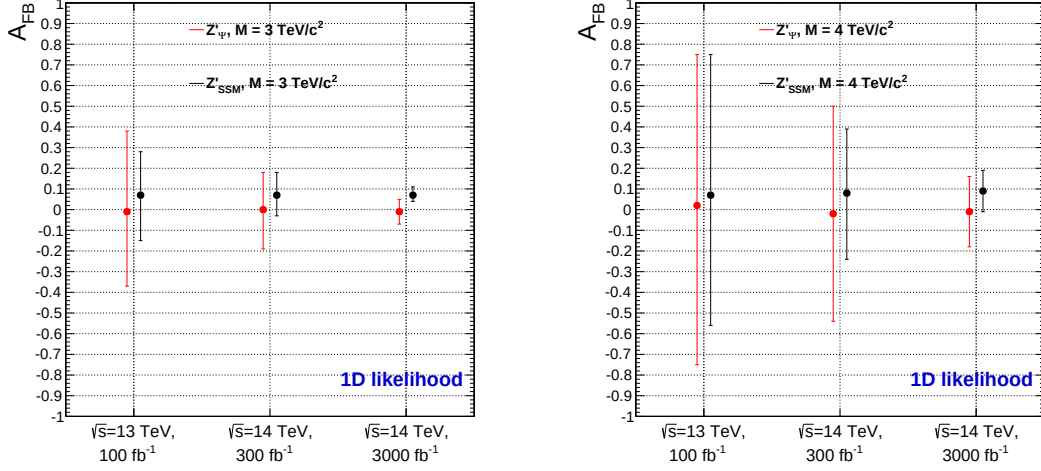


Figure 7.16: Expected median and 1σ range on the A_{FB} measurement for a Z'_ψ and Z'_{SSM} resonance at 3 or 4 TeV for various integrated luminosities. The results presented here only use the $\cos\theta_{CS_{meas}}$ information.

as a function of A_{FB} for a pseudo-experiment corresponding to a Z'_ψ signal at 3 TeV with 3000 fb^{-1} of integrated luminosity. The distribution of the estimate of A_{FB} for 1000 pseudo-experiment is shown in figure 7.19 and a fit to a Gaussian is performed. The width of the Gaussian is found to be reduced by around 10% when the 2D likelihood is used.

7.3.3 gg - $q\bar{q}$ fractions

For a spin 0 or 2 resonances, both gg and $q\bar{q}$ production modes are possible. For a given proton-proton center of mass energy, these fractions depend on the resonance mass. In this section the measurement of the gg and $q\bar{q}$ fractions for a G_{RS} is investigated.

7.3.3.1 Procedure

The procedure is the same as the one described in section 7.3.2 for the A_{FB} measurement. One first builds the pdfs p_{gg} and $p_{q\bar{q}}$ corresponding to the two extreme cases (100% gg or 100% $q\bar{q}$) and looks for κ that minimizes the likelihood given by:

$$p_\kappa = \kappa \times p_{gg} + (1 - \kappa) \times p_{q\bar{q}} \quad (7.12)$$

Here κ is directly the gg fraction.

7.3.3.2 Results

The results for a G_{RS} are shown in tables 7.10 and 7.11 for the 1D and 2D likelihood. As expected, the 2D likelihood slightly improves the precision of the measurement. The theoretical values are listed in table 7.12. The 2D distributions for a graviton purely

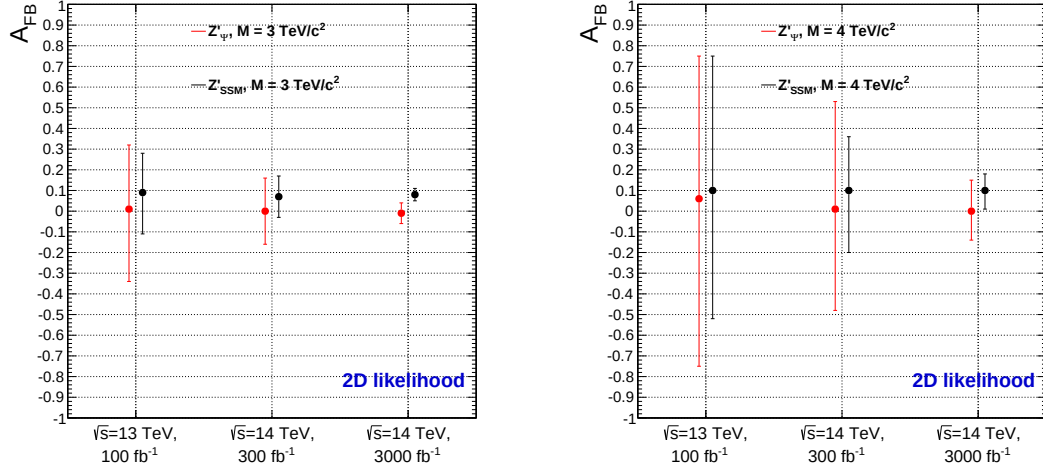


Figure 7.17: Expected median and 1σ range on the A_{FB} measurement for a Z'_ψ and Z'_{SSM} resonance at 3 or 4 TeV for various integrated luminosities. The results presented here use the $\cos\theta_{CS_{meas}}$ and y_{ee} information.

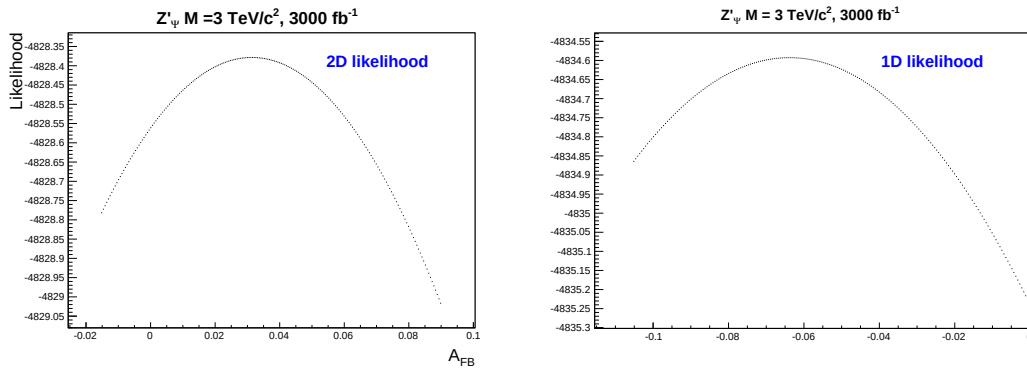


Figure 7.18: Likelihood versus A_{FB} curve obtained for one random pseudo-experiment corresponding to a Z'_ψ signal at 3 TeV with 3000 fb^{-1} of integrated luminosity at $\sqrt{s} = 14 \text{ TeV}$. The likelihood is calculated from a 1D (left) or 2D (right) pdf.

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Model	100 fb ⁻¹ (13 TeV)	300 fb ⁻¹ (14 TeV)	3000 fb ⁻¹ (14 TeV)
Z'_Ψ (M = 3 TeV/c ²) 1D	$-0.01^{+0.39}_{-0.36}$	$0.00^{+0.18}_{-0.19}$	$-0.01^{+0.06}_{-0.06}$
Z'_{SSM} (M = 3 TeV/c ²) 1D	$0.07^{+0.21}_{-0.22}$	$0.07^{+0.11}_{-0.10}$	$0.07^{+0.04}_{-0.03}$
Z'_Ψ (M = 4 TeV/c ²) 1D	$0.02^{+0.73}_{-0.77}$	$-0.02^{+0.52}_{-0.52}$	$-0.01^{+0.17}_{-0.17}$
Z'_{SSM} (M = 4 TeV/c ²) 1D	$0.07^{+0.68}_{-0.63}$	$0.08^{+0.31}_{-0.32}$	$0.09^{+0.10}_{-0.10}$
DY (1056 evts) M > 3 TeV/c ² 1D	$0.59^{+0.06}_{-0.06}$		

Table 7.7: Expected median and 1σ range on the A_{FB} measurement for a Z'_Ψ and Z'_{SSM} resonance at 3 or 4 TeV for various integrated luminosities. The results are also presented for 1056 Drell-Yan events with M>3 TeV. The results presented here only use the $\cos\theta_{CS_{meas}}$ information.

Model	100 fb ⁻¹ (13 TeV)	300 fb ⁻¹ (14 TeV)	3000 fb ⁻¹ (14 TeV)
Z'_Ψ (M = 3 TeV/c ²) 2D	$0.01^{+0.31}_{-0.35}$	$0.00^{+0.16}_{-0.16}$	$-0.01^{+0.05}_{-0.05}$
Z'_{SSM} (M = 3 TeV/c ²) 2D	$0.09^{+0.19}_{-0.20}$	$0.07^{+0.10}_{-0.10}$	$0.08^{+0.03}_{-0.03}$
Z'_Ψ (M = 4 TeV/c ²) 2D	$0.06^{+0.69}_{-0.81}$	$0.01^{+0.52}_{-0.49}$	$-0.00^{+0.15}_{-0.14}$
Z'_{SSM} (M = 4 TeV/c ²) 2D	$0.10^{+0.65}_{-0.62}$	$0.10^{+0.26}_{-0.30}$	$0.10^{+0.08}_{-0.09}$
DY (1056 evts) M > 3 TeV/c ² 2D	$0.61^{+0.06}_{-0.06}$		

Table 7.8: Expected median and 1σ range on the A_{FB} measurement for a Z'_Ψ and Z'_{SSM} resonance at 3 or 4 TeV for various integrated luminosities. The results are also presented for 1056 Drell-Yan events with M>3 TeV. The results presented here use the $\cos\theta_{CS_{meas}}$ and y_{ee} information.

	up quarks	down quarks	$\sqrt{s} = 13$ TeV pp collisions
Z'_Ψ	0	0	0
Z'_{SSM}	0.075	0.105	0.08
Drell-Yan M > 3 TeV	0.595	0.625	0.60

Table 7.9: Theoretical A_{FB} at leading order for different $q\bar{q} \rightarrow l\bar{l}$ processes. Because of their different couplings, up and down quarks induce in general different asymmetries. They are therefore presented separately, as well as the total asymmetry for pp collisions at $\sqrt{s} = 13$ TeV. The values quoted for the signal samples do not take interference effects with the Drell-Yan process into account.

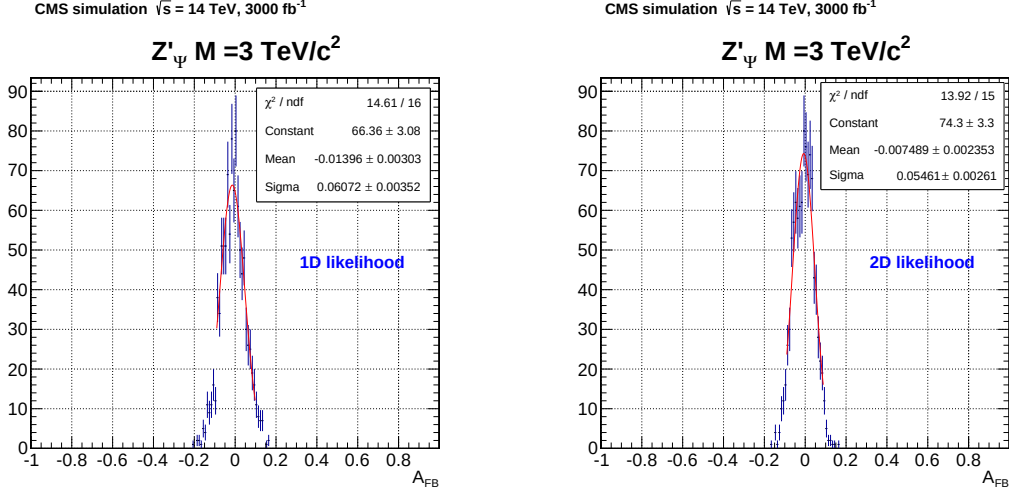


Figure 7.19: Distributions of the A_{FB} estimate for 1000 pseudo-experiments corresponding to a Z'_ψ signal at 3 TeV with 3000 fb^{-1} of integrated luminosity at $\sqrt{s} = 14 \text{ TeV}$. The likelihood is calculated from a 1D (left) or 2D (right) pdf. A fit to a gaussian is performed. The width is reduced by $\approx 10\%$ when using the 2D pdf.

Model	100 fb ⁻¹ (13 TeV)	300 fb ⁻¹ (14 TeV)	3000 fb ⁻¹ (14 TeV)
G_{RS} ($M = 3 \text{ TeV}/c^2$) 1D	$0.35^{+0.33}_{-0.32}$	$0.36^{+0.16}_{-0.15}$	$0.36^{+0.05}_{-0.05}$
G_{RS} ($M = 4 \text{ TeV}/c^2$) 1D	$0.34^{+0.67}_{-0.34}$	$0.19^{+0.56}_{-0.19}$	$0.24^{+0.15}_{-0.16}$

Table 7.10: Expected median and 1σ range on the gg fraction measurement for a G_{RS} resonance at 3 or 4 TeV for various integrated luminosities. The results presented here only use the $\cos\theta_{CS_{meas}}$ information.

produced through gg or $q\bar{q}$ fusions are shown in figure 7.20. gg fusion favours low y_{ee} values whereas $q\bar{q}$ favours high y_{ee} values. The drop of events at $|\cos\theta_{CS}| \approx 0.6$ for large $|y_{ee}|$ in the $q\bar{q}$ case is an acceptance effect (one of the electrons is often emitted close to or in the transition region between the ECAL barrel and endcaps).

7.4 Summary

One of the straightforward ways to observe physics beyond the Standard Model at the LHC would be through the emergence of a new resonance in the dielectron mass spectrum at the TeV scale. Although the search for a signal in the data collected by CMS during the LHC Run 1 and detailed in chapter 6 did not lead to such a discovery, a new signal could possibly show up only a couple of months after the LHC restarts in 2015 thanks to the increase of the LHC beam energy. In case a signal is observed, the leptons angular distribution provide valuable information to distinguish between different new physics

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Model	100 fb ⁻¹ (13 TeV)	300 fb ⁻¹ (14 TeV)	3000 fb ⁻¹ (14 TeV)
G_{RS} (M = 3 TeV/c ²) 2D	0.37 ^{+0.31} _{-0.30}	0.36 ^{+0.13} _{-0.13}	0.37 ^{+0.05} _{-0.04}
G_{RS} (M = 4 TeV/c ²) 2D	0.34 ^{+0.66} _{-0.34}	0.20 ^{+0.45} _{-0.20}	0.24 ^{+0.13} _{-0.15}

Table 7.11: Expected median and 1σ range on the gg fraction measurement for a G_{RS} resonance at 3 or 4 TeV for various integrated luminosities. The results presented here use the $\cos\theta_{CS_{meas}}$ and y_{ee} information.

RS graviton M = 3 TeV	0.38
RS graviton M = 4 TeV	0.25

Table 7.12: Theoretical gg fractions at leading order for a Randall-Sundrum resonance at 3 or 4 TeV.

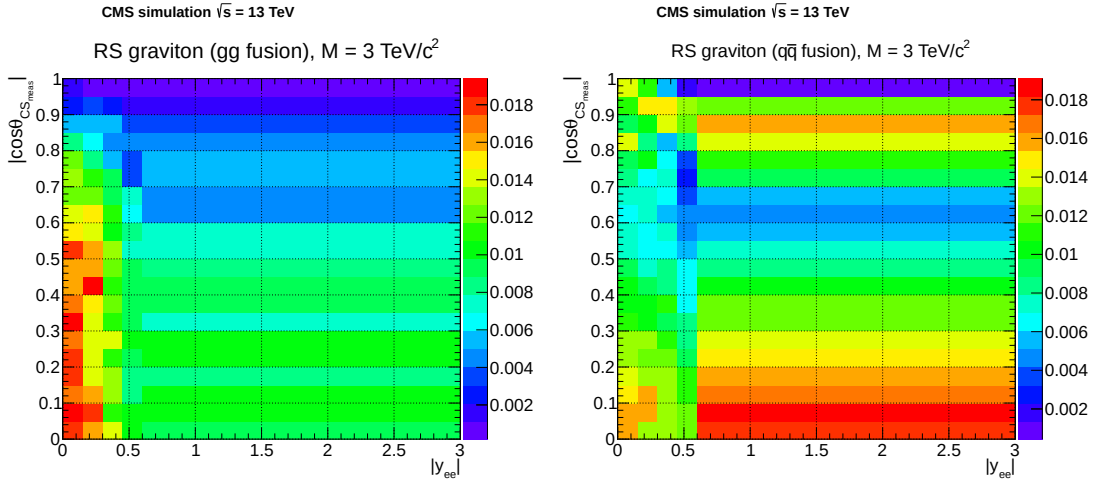


Figure 7.20: $(\cos\theta_{CS_{meas}}, |y_{ee}|)$ distributions for a G_{RS} purely produced through gg (left) or $q\bar{q}$ (right) fusion. The average rapidity is greater in the second case. The collisions are indeed more forward for $q\bar{q}$ pairs, for which the quark carries in average more momentum than the antiquark. The drop of events at $|\cos\theta_{CS}| \approx 0.6$ for large $|y_{ee}|$ in the $q\bar{q}$ case is an acceptance effect (one of the electrons is often emitted close to or in the gap).

models and to measure the spin, and other properties of the observed signal. In this chapter, the precision that could be reached in the measurement of the spin, forward-backward asymmetry and production modes of the new signal were studied.

Three resonant signal models were considered: the sequential Standard Model Z'_{SSM} , the Z'_Ψ (spin 1) and the G_{RS} (spin 2). For a Z'_Ψ or a G_{RS} with a mass of 3 TeV, the wrong spin hypotheses can be excluded by more than 2σ for $\sqrt{s} = 14$ TeV and 300 fb^{-1} of integrated luminosity. An exclusion at the 5σ level requires however 3000 fb^{-1} of integrated luminosity. For a 4 TeV signal, the full phase 2 luminosity is needed to reach the $2\text{-}3\sigma$ level. Concerning the forward-backward asymmetry measurement, 300 fb^{-1} would lead to a measurement of 0.00 ± 0.18 for a 3 TeV Z'_Ψ , to be compared with the Drell-Yan expected value of 0.60. A factor 3 smaller is reached for 3000 fb^{-1} . As for the spin case, the full phase 2 luminosity is needed to make any statement about the forward-backward asymmetry for 4 TeV signals.

These results are extracted mainly from the shape of the $\cos\theta_{CS}$ distribution. It is however shown that a sizeable improvement of the precision of the measurements can be achieved by using the dielectron pair rapidity information.

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Conclusions

The start-up of the LHC in 2010 at CERN has marked the beginning of a new era in particle physics. Thanks to the high energy and intensity of its proton beams, it is possible to study or search for very rare processes in an energy range that reaches now the TeV.

In three years of activity, the two main experiments at the LHC, the ATLAS and CMS experiments, achieved their first breakthrough by confirming the existence of the so far only hypothetical scalar boson, postulated by the electroweak symmetry breaking mechanism. The discovery of this boson represents a huge success for the current theory describing the elementary particles and their interactions at the subatomic level, the Standard Model. Despite its predictions experimentally verified with great precision, it is broadly believed that the Standard Model only consists of the low energy residual of a more fundamental theory which could for instance explain why gravity appears so weak with respect to the other interactions at subatomic distances. Several theoretical considerations suggest that this new physics could be characterized by an energy in the TeV range.

This thesis presented the results of a search with the CMS detector for new narrow heavy (>200 GeV) resonances decaying into dielectron pairs. The invariant mass spectrum of dielectron pairs is measured and a bump hunting is performed. This search takes full advantage of the high energy of the LHC proton beams and the low background from Standard Model in the signal region, together with the good electron energy resolution (around 1% for electron energies above 100 GeV), makes the signature studied a smoking gun for new physics.

In chapter 1, after briefly introducing the Standard Model, we highlighted some of its shortcomings and reviewed various models of new physics introducing new heavy resonances decaying into electron-positron pair. Examples are the Z'_ψ , a new heavy gauge boson predicted in some theories where the three interactions of the Standard Model are unified at high energy, and the Randall-Sundrum graviton, G_{RS} , which arises in models where the existence of extra spatial dimensions is postulated to explain the apparent weakness of the gravity.

Chapter 2 was devoted to the description of the experimental apparatus. We presented the data taking conditions during the first years of the LHC operation and the projections for the future. After delivering around 5 fb^{-1} of integrated luminosity at a center of mass energy ($\sqrt{s}$) of 7 TeV in 2011 and 20 fb^{-1} at $\sqrt{s} = 8$ TeV in 2012, the LHC is currently in a shutdown period. It is expected to resume soon (early 2015) and will generate proton-proton collisions with an increased $\sqrt{s}$ of 13 TeV. A long term upgrade of the LHC and of the detector (LHC phase 2) is also considered, in order to reach instantaneous luminosities 5 times higher than the nominal value of $10^{-34} \text{ cm}^{-2}\text{s}^{-1}$. The various components forming

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the CMS detector were also presented in chapter 2 and the general reconstruction of the physical objects produced during a collision (electrons, photons, jets resulting from the hadronisation of a quark or a gluon,...) were described.

Chapter 3 explained the current modelling of a hard proton-proton collision, described as an interaction between one parton from each proton, and the various successive steps followed for their simulation and reconstruction with the detector. It was stressed that the production cross section of parton pairs with large fraction of the proton momentum is extremely dependent of $\sqrt{s}$.

One of the key points of the search for heavy dielectron resonance is an efficient and robust selection of very high energy electrons, with energies possibly above 1 TeV. This was the topic of chapter 4. The basic electron reconstruction requires the association of a track reconstructed in the tracker and an energy deposit in the electromagnetic calorimeter. In hadron colliders, cross sections involving jets are by several orders of magnitude higher than those involving electrons, and jets incorrectly reconstructed as electrons in the detector constitute a non negligible experimental background. We therefore developed and presented a serie of cuts designed to reject those fakes while remaining as efficient as possible for electrons, paying a special attention to high energy electrons. One of the main criteria used to distinguish an electron from a jet is the calorimeter isolation: no nearby energy activity is expected around the electron and we therefore applied a cut on the total transverse energy collected in the hadronic and electromagnetic calorimeters in a cone centered on the electron candidate direction. We however showed that the presence of pile-up collisions can highly impact the isolation, bringing the efficiency of this cut down to 0% when 50 pile-up collisions are recorded, for electrons with a transverse momentum of 35 GeV. A correction was then applied to subtract this pile-up contribution from the isolation variable allowing us to recover a high and pile-up independent efficiency. The extension of this correction to other variables was also studied.

The electron selection discussed in chapter 4 was designed using simulations. Its efficiency directly affects the yield of a possible signal and must be estimated as precisely as possible. It is therefore crucial to check that the simulations are well describing the data and that no unexpected drop happens at high energy. The electrons in the signal are however characterized by a transverse energy E_T up to 1 TeV, and the low statistics available in that region in the data makes a measurement impossible. There is no other choice than using the simulation estimate of the efficiency there. At lower E_T however, the Drell-Yan process ($q\bar{q} \rightarrow l\bar{l}$) provides an abundant statistics of electrons in the data. The strategy adopted is therefore to measure a scale factor, defined as the ratio of the electron selection efficiency measured on Drell-Yan events in both data and simulations and to study its evolution from moderate E_T to E_T as high as possible. Various methods were developed in chapter 4 to measure the scale factor. They both gave consistant results. The scale factors measured for electrons reconstructed in the barrel were found to be independent of E_T and compatible with 0.98 within 2% for E_T up to 500 GeV. A constant scale factor with a 2% systematic uncertainty on its evolution was therefore assumed. The measurement method for electrons reconstructed in one of the endcaps were similar but the reduced statistics at very high E_T in this channel led us to assign a 4% systematic uncertainty on the scale factor evolution in the endcaps. In addition to the selection efficiency, we also measured efficiencies for the trigger and reconstruction

steps. They both turned out to be consistent and very close to 1 resulting in scale factors compatible with 1 and with an uncertainty below the one obtained for the selection efficiency.

Chapter 6 detailed the results of the search for new heavy dielectron resonances and described the various steps of the analysis. The event selection, which only requires two electrons whose at least one is reconstructed in the barrel, is deliberately simple to keep it as efficient and as model independent as possible. For a resonant spin 1 signal above 1 TeV, the total acceptance times efficiency is 65%. The dielectron mass resolution was estimated from simulations but was corrected by the observed difference between data and simulations at for dielectron pairs at the Z peak. The main Standard Model background comes from the Drell-Yan process and is irreducible. It is estimated from simulations. The next background contribution comes from leptonic decays of pairs of heavy bosons or of a $t\bar{t}$ pair. It is estimated from simulations but a cross check is performed on the data exploiting the fact that those processes represent the leading contribution of the $e-\mu$ pair invariant mass spectrum. The dielectron invariant mass distribution is then measured and no significant excess above the background expectation is observed. Upper limits at 95% confidence level are therefore set on R_σ , defined as the ratio of the production cross section of a new narrow heavy resonance to the production cross section of the Drell-Yan process in the Z peak region ($60 < M_{ee}(\text{GeV}) < 120$). A Bayesian approach, using an unbinned likelihood based on the dielectron pair invariant mass distribution, is followed. This allows us to fully exploit the difference in shape between the signal and the background. The obtained limits are generic and hold for any signal whose intrinsic width is smaller or comparable to the experimental resolution on the dielectron pair mass. Three different acceptance times efficiency associated to three spin scenarios for the resonance are considered. Similar limits are obtained in the three cases at high mass and a upper maximal ratio of $R_\sigma < 2.10^{-7}$ is reached for a resonance with a mass above 2 TeV. The upper limits can be reinterpreted as a lower mass condition on resonances predicted by new physics models, such as the two described above. Limits for a new boson with couplings identical to those of the Z boson (Z'_{SSM}) can also be calculated. The presented analysis excludes a Z'_ψ or a Z'_{SSM} below 2.34 TeV or 2.67 TeV respectively. Similarly, for the Randall-Sundrum graviton with a coupling $c=0.05$ ($c=0.1$), a lower limit of 2.13 TeV (2.50 TeV) is set.

Despite the absence of signal so far, it is possible that a signal emerges fastly in the next data taking at $\sqrt{s} = 13$ TeV. The production cross section of objects with a mass in the TeV range is indeed very sensitive to $\sqrt{s}$. Chapter 7 quantified, assuming such a signal is found, the precision that could be reached in the measurement of the spin, the forward-backward asymmetry and the production modes of a newly discovered dielectron resonance. The measurements are based on the study of the dielectron angular distributions and on their rapidity. For the three signals considered at 3 TeV, the wrong spin hypotheses can be excluded by more than 2σ for $\sqrt{s} = 14$ TeV and 300 fb^{-1} of integrated luminosity. An exclusion at the 5σ level requires however 3000 fb^{-1} of integrated luminosity. For a 4 TeV signal, the full phase 2 luminosity is needed to reach the $2\text{-}3\sigma$ level. Concerning the forward-backward asymmetry measurement, 300 fb^{-1} would lead to a measurement of 0.00 ± 0.18 for a 3 TeV Z'_ψ , to be compared with the Drell-Yan expected value of 0.60. A factor 3 smaller is reached for 3000 fb^{-1} . As for the spin case,

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the full phase 2 luminosity is needed to make any statement about the forward-backward asymmetry for 4 TeV signals.

The hunt for new physics at the LHC involves many physicists and the search presented in this thesis consists of only one of the many searches carried out by the CMS collaboration. The discovery of new physics might need the combination of the analyses of several final states and topologies. In the search for a massive resonance, the dielectron and dimuon channels are already combined. The combination with other channels, in particular the diphoton channel, should allow us to improve the sensitivity to the production of new particles with similar couplings to all the Standard Model particles, such as gravitons. Another possible improvement of the analysis is the consideration of other variables in addition to the dielectron mass, such as the forward-backward asymmetry.

Despite an intense effort to track it, no hint of new physics was observed at the LHC so far. As already advertised though, the discovery potential of the LHC will strongly extend in the next year. If the first years of the LHC will enter the history of particle physics as the discovery years of the scalar boson, the upcoming ones might be the ones of new physics.

Appendix A

Drell-Yan cross section calculation

A.1 Conventions

In this appendix, the following conventions are used:

- The metric of the Minkowsky space $g^{\mu\nu}$ is defined with $sign(g) = (1, -1, -1, -1)$
- 4-dimensional vectors are represented as simple characters (v)
- 3-dimensional vectors are represented with an arrow ($\vec{v}$)
- Repeated greek indices in an equation implies a sum ($v^\mu w_\mu = v^0 w_0 + v^1 w_1 + v^1 w_1 + v^1 w_1$)
- The $(. .)$ syntax denotes a Minkoswki product ($(v.w) = v^\mu w_\mu$)
- Unless explicately stated the unit system is chosen such that $\hbar = c = \epsilon_0 = 1$

A.2 Introduction and hypotheses

The Drell-Yan process is defined as the annihilation of a quark-antiquark pair into a lepton-antilepton pair. This process is described at leading order by the two Feynman diagrams of figure 1.2. These two amplitudes are proportional to the fine structure constant $\alpha \approx 1/137$.

In the Standard Model, there is actually a third diagram contributing, namely the production of the scalar boson (H^0) predicted by the electroweak symmetry breaking mechanism.

However, the fermion coupling to the scalar boson is proportional to m_f/v where m_f is the mass of the fermion and v is the vacuum expected value of the scalar field and is measured to be 246 GeV in the Standard Model. This leads to an amplitude proportional to $\frac{m_q(GeV)}{246} \times \frac{m_l(GeV)}{246}$ for this diagram. The heaviest lepton, the tau, has a mass of ≈ 1.8 GeV while the second heaviest quark is the bottom with a mass of ≈ 4.7 GeV. For all these particles, this amplitude is therefore small compared to the previous ones. As quarks are never free in the Nature, they are experimentally produced during hadrons collisions.

A. DRELL-YAN CROSS SECTION CALCULATION

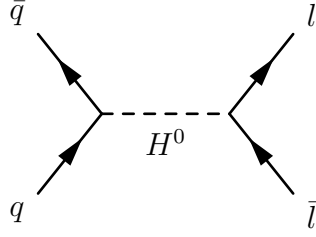


Figure A.1: Feynman diagram of the process $q + \bar{q} \rightarrow H^0 \rightarrow l + \bar{l}$ at leading order.

The production rate of each quark depends of its parton density inside the hadrons. As a consequence:

- The contribution of the proton valence implies that up and down quarks ($m < 10$ MeV) are produced in larger amounts than the others.
- Heavier quarks are harder produced because of their mass.
- A top-antitop initial state cannot occur below an invariant mass of $2 \times m_{top} \approx 350$ GeV.

In these conditions, the two first diagrams are dominant by several orders of magnitude and the latter will neglected in the calculation. Finally, only Drell-Yan events with an energy in the center of mass corresponding to the Z mass (91.19 GeV) or higher are studied in this thesis. This allows one to work in the ultra-relativistic limit and neglect the quarks and leptons masses.

A.3 Cross section and forward-backward asymmetry formula

The differential cross section of the scattering $\mathcal{A} + \mathcal{B} \rightarrow 1 + 2 + \dots + n$ is given by:

$$d\sigma_{\mathcal{A}+\mathcal{B} \rightarrow 1+2+\dots+n} = \frac{1}{2E_{\mathcal{A}}2E_{\mathcal{B}}|v_{\mathcal{A}} - v_{\mathcal{B}}|} \cdot (2\pi)^4 \delta^{(4)} \left(p_{\mathcal{A}} + p_{\mathcal{B}} - \sum_{f=1,\dots,n} p_f \right) \cdot \prod_{f=1,\dots,n} \left(\frac{d^3\vec{p}_f}{(2\pi)^3 2E_f} \right) \cdot X \quad (\text{A.1})$$

where p_i , $\vec{p}_i$, E_i are respectively the four-momentum, the three-momentum and the energy of the particle i and $|v_{\mathcal{A}} - v_{\mathcal{B}}|$ is the difference in velocity between $\mathcal{A}$ and $\mathcal{B}$ and can be expressed as:

$$|v_{\mathcal{A}} - v_{\mathcal{B}}| = \frac{1}{E_{\mathcal{A}}E_{\mathcal{B}}} \sqrt{(p_{\mathcal{A}} \cdot p_{\mathcal{B}})^2 - (m_{\mathcal{A}}m_{\mathcal{B}})^2} \quad (\text{A.2})$$

Finally, X contains the matrix element:

$$X = |\mathcal{M}_{\mathcal{A}+\mathcal{B} \rightarrow 1+2+\dots+n}|^2 \quad (\text{A.3})$$

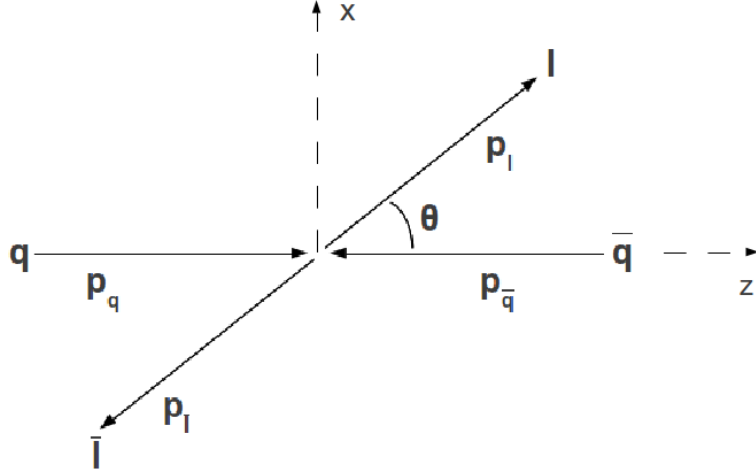


Figure A.2: Definition of the frame chosen for the calculation. The axes are chosen such that the z-axis points to the quark direction of flight and that the y-axis is perpendicular to the plan defined by the quarks and the leptons.

The forward-backward asymmetry is defined in the $q\bar{q}$ center of mass frame and all the calculations will therefore be made in this frame. The axes are chosen such that the z-axis points to the quark direction of flight and that the y-axis is perpendicular to the plan defined by the quark and the lepton directions. The angle θ is defined as the angle between the directions of flight of the quark and the lepton. The definition of the frame is illustrated in figure A.3. The forward-backward asymmetry is obtained starting from the forward and backward cross sections given by $\sigma_F = \sigma_{\theta < \pi/2}$ and $\sigma_B = \sigma_{\theta > \pi/2}$. The forward-backward asymmetry is then defined by:

$$A_{FB} = \frac{\sigma_F - \sigma_B}{\sigma_F + \sigma_B} \quad (\text{A.4})$$

The two extreme cases (+1 and -1) correspond to a purely forward or backward cross section (the lepton is always closer to the quark or to the antiquark).

In the ultrarelativistic limit, calling $\sqrt{s'}$ the total energy, one has $E_q = E_{\bar{q}} = E_l = E_{\bar{l}} = \sqrt{s'}/2$ and therefore:

$$p_q = \left(\frac{\sqrt{s'}}{2}, 0, 0, \frac{\sqrt{s'}}{2} \right) \quad (\text{A.5})$$

$$p_{\bar{q}} = \left(\frac{\sqrt{s'}}{2}, 0, 0, -\frac{\sqrt{s'}}{2} \right) \quad (\text{A.6})$$

$$p_l = \left(\frac{\sqrt{s'}}{2}, \frac{\sqrt{s'}}{2} \sin \theta, 0, \frac{\sqrt{s'}}{2} \cos \theta \right) \quad (\text{A.7})$$

$$p_{\bar{l}} = \left(\frac{\sqrt{s'}}{2}, -\frac{\sqrt{s'}}{2} \sin \theta, 0, -\frac{\sqrt{s'}}{2} \cos \theta \right) \quad (\text{A.8})$$

The cross section (equation A.1) for the Drell-Yan process in this frame is:

A. DRELL-YAN CROSS SECTION CALCULATION

$$\begin{aligned}
d\sigma_{q+\bar{q} \rightarrow l+\bar{l}} &= \frac{1}{2E_q 2E_{\bar{q}} |v_q - v_{\bar{q}}|} (2\pi)^4 \delta^{(4)}(p_q + p_{\bar{q}} - p_l - p_{\bar{l}}) \\
&\quad \times \frac{d^3 \vec{p}_l}{(2\pi)^3 2E_l} \frac{d^3 \vec{p}_{\bar{l}}}{(2\pi)^3 2E_{\bar{l}}} \cdot X \\
&= \frac{1}{8\pi^2 s'^2} \cdot \delta^{(4)}(p_q + p_{\bar{q}} - p_l - p_{\bar{l}}) \cdot d^3 \vec{p}_l d^3 \vec{p}_{\bar{l}} \cdot X
\end{aligned} \tag{A.9}$$

The last equation uses the fact that from A.2: $|v_q - v_{\bar{q}}| = 2$. Using the following change of variables:

$$\begin{cases} u = p_l + p_{\bar{l}} \\ v = p_l - p_{\bar{l}} \end{cases} \tag{A.10}$$

one has:

$$\delta^{(4)}(p_q + p_{\bar{q}} - p_l - p_{\bar{l}}) d^3 \vec{p}_l d^3 \vec{p}_{\bar{l}} = \frac{1}{8} \delta^{(4)}(u - v) d^3 \vec{u} d^3 \vec{v} \tag{A.11}$$

(the $1/8$ factor comes from the Jacobian). The delta function implies that the integral over $\vec{u}$ can be safely performed without worrying about the content of X . Furthermore $d^3 \vec{v}$ can be reexpressed as $|\vec{v}|^2 d|\vec{v}| d\Omega$ where $d\Omega$ is a solid angle element. One is left with $\delta^{(0)}(u^0 - \sqrt{s'})$. But $u^0 = 2E_l = 2|\vec{p}_l| = |\vec{p}_l - \vec{p}_{\bar{l}}| = |\vec{v}|$ and the integral on $|\vec{v}|$ can be performed too. One then obtains the value of the cross section per unit of solid angle:

$$\frac{d\sigma}{d\Omega} = \frac{1}{64\pi^2 s'} \cdot X \tag{A.12}$$

The last thing to compute is X . In the Drell-Yan process, there are two fermions in the initial and final states with unknown polarization. The quarks and antiquarks exist in three different colors but only colorless pairs can form a photon or a Z boson. To compute X , one must therefore sum the amplitude corresponding to any possible spin configurations, average on the spins of the incoming particles and require the quark and antiquark to carry an opposite color. This leads to:

$$X = \frac{1}{3} \frac{1}{2} \sum_{s_q} \frac{1}{2} \sum_{s_{\bar{q}}} \sum_{s_l} \sum_{s_{\bar{l}}} |\mathcal{M}_{s_q, s_{\bar{q}}, s_l, s_{\bar{l}}}(q + \bar{q} \rightarrow l + \bar{l})|^2 \tag{A.13}$$

The Feynman rules presented in section A.7 group the 2 helicity states together using $\psi = \psi_L + \psi_R$ reducing the calculation to only one amplitude per intermediate boson.

A.4 Calculation of the matrix elements

A.4.1 $\mathcal{M}_\gamma$

From the Feynman rules for the Standard Model, the matrix element is given at leading order by:

$$i\mathcal{M}_\gamma = \frac{iQ_l Q_q e^2}{s'} \bar{u}(p_l, s_l) \gamma^\mu v(p_{\bar{l}}, s_{\bar{l}}) \bar{v}(p_{\bar{q}}, s_{\bar{q}}) \gamma_\mu u(p_q, s_q) \tag{A.14}$$

A.4.2 $\mathcal{M}_Z$

The same calculation for the Z boson matrix element leads to:

$$i\mathcal{M}_Z = \frac{ig^2}{(2\cos\theta_W)^2} \frac{\bar{u}(p_l, s_l)\gamma^\mu(g_{V_l} + g_{A_l}\gamma^5)v(p_{\bar{l}}, s_{\bar{l}})\bar{v}(p_{\bar{q}}, s_{\bar{q}})\gamma_\mu(g_{V_q} + g_{A_q}\gamma^5)u(p_q, s_q)}{s' - M_Z + is'\Gamma_Z/M_Z} \quad (\text{A.15})$$

$$= \frac{iQ_l Q_q e^2}{s'} \mathcal{R} \bar{u}(p_l, s_l)\gamma^\mu(g_{V_l} + g_{A_l}\gamma^5)v(p_{\bar{l}}, s_{\bar{l}})\bar{v}(p_{\bar{q}}, s_{\bar{q}})\gamma_\mu(g_{V_q} + g_{A_q}\gamma^5)u(p_q, s_q) \quad (\text{A.16})$$

where $e = g \sin\theta_W$ and the following convenient notations are defined:

$$\mathcal{R} = \frac{1}{Q_l Q_q \sin^2 2\theta_W} \frac{s'}{s' - M_Z^2 + is'\Gamma_Z/M_Z} \quad (\text{A.17})$$

$$g_{V_{l,q}} = I_{Wl,q}^3 - 2Q_{l,q} \sin^2 \theta_W \quad (\text{A.18})$$

$$g_{A_{l,q}} = -I_{Wl,q}^3 \quad (\text{A.19})$$

M_Z and Γ_Z are respectively the Z boson mass and total decay width. To save some papers and avoid equation overflow, one sets $\psi(p_i, s_i) \equiv \psi_i$ in the next sections.

A.5 Squaring the matrix elements

The total cross section is proportional to the square of the sum of the different amplitudes. There are therefore three contributions to be computed: $|\mathcal{M}_\gamma|^2$, $|\mathcal{M}_Z|^2$ and the interference $\mathcal{M}_\gamma \mathcal{M}_Z^* + c.c.$

A.5.1 $|\mathcal{M}_\gamma|^2$

The complex conjugate of $\mathcal{M}_\gamma$ is:

$$\mathcal{M}_\gamma^* = \frac{Q_l Q_q e^2}{s'} \bar{v}_{\bar{l}} \gamma^\nu u_{\bar{q}} \bar{u}_l \gamma_\nu v_{\bar{q}} \quad (\text{A.20})$$

where one used the fact that $(\bar{\psi}\gamma^\mu\chi)^* = (\psi^\dagger\gamma^0\gamma^\mu\chi)^* = \chi^\dagger(\gamma^\mu)^\dagger(\gamma^0)^\dagger\psi = \chi^\dagger\gamma^0\gamma^\mu\gamma^0\gamma^0\psi = \bar{\chi}\gamma^\mu\psi$ with the help of the relations A.60, A.63. This product is a complex number and can therefore be safely moved inside an equation. The matrix element averaged over the incoming particles spins is then:

$$\frac{1}{4} \sum_{s_q s_{\bar{q}} s_l s_{\bar{l}}} |\mathcal{M}_\gamma|^2 = \frac{Q_l^2 Q_q^2 e^4}{4s'^2} \sum_{s_q s_{\bar{q}} s_l s_{\bar{l}}} (\bar{u}_l \gamma^\mu v_{\bar{l}} \bar{v}_{\bar{l}} \gamma^\nu u_l) (\bar{v}_{\bar{q}} \gamma_\mu u_q \bar{u}_q \gamma_\nu v_{\bar{q}}) \quad (\text{A.21})$$

In the first brackets term, one has, writing down explicitly the sum over the spinor indices and using equations A.68 and A.69:

$$\sum_{s_l s_{\bar{l}}} \bar{u}_l \gamma^\mu v_{\bar{l}} \bar{v}_{\bar{l}} \gamma^\nu u_l = \sum_{s_l s_{\bar{l}}} \sum_{\alpha\beta\gamma\delta} \bar{u}_{l\alpha} (\gamma^\mu)_{\alpha\beta} v_{\bar{l}\beta} \bar{v}_{\bar{l}\gamma} (\gamma^\nu)_{\gamma\delta} u_{l\delta} \quad (\text{A.22})$$

$$= \sum_{\alpha\beta\gamma\delta} (p_{l\rho} \gamma^\rho + m_l)_{\delta\alpha} (p_{\bar{l}\sigma} \gamma^\sigma - m_{\bar{l}})_{\beta\gamma} (\gamma^\mu)_{\alpha\beta} (\gamma^\nu)_{\gamma\delta} \quad (\text{A.23})$$

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In the ultrarelativistic limit, it comes, using equation A.65:

$$\sum_{s_l s_{\bar{l}}} \bar{u}_l \gamma^\mu v_{\bar{l}} \bar{v}_l \gamma^\nu u_l = \text{Tr} (p_{l_\rho} p_{\bar{l}_\sigma} \gamma^\rho \gamma^\mu \gamma^\sigma \gamma^\nu) \quad (\text{A.24})$$

$$= 4 p_{l_\rho} p_{\bar{l}_\sigma} (g^{\rho\mu} g^{\sigma\nu} - g^{\rho\sigma} g^{\mu\nu} + g^{\rho\nu} g^{\mu\sigma}) \quad (\text{A.25})$$

$$= 4 (p_l^\mu p_{\bar{l}}^\nu - (p_l \cdot p_{\bar{l}}) g^{\mu\nu} + p_l^\nu p_{\bar{l}}^\mu) \quad (\text{A.26})$$

The same procedure can be applied for the quark part (second brackets) of equation A.21. Putting all together one finds:

$$\begin{aligned} \frac{1}{4} \sum_{s_q s_{\bar{q}} s_l s_{\bar{l}}} |\mathcal{M}_\gamma|^2 &= \frac{Q_l^2 Q_q^2 e^4}{4s'^2} 16 \times (p_l^\mu p_{\bar{l}}^\nu - (p_l \cdot p_{\bar{l}}) g^{\mu\nu} + p_l^\nu p_{\bar{l}}^\mu) (p_{q_\mu} p_{\bar{q}_\nu} - (p_q \cdot p_{\bar{q}}) g_{\mu\nu} + p_{q_\mu} p_{\bar{q}_\nu}) \\ &= \frac{Q_l^2 Q_q^2 e^4}{4s'^2} 16 \times 2 [(p_l \cdot p_q)(p_{\bar{l}} \cdot p_{\bar{q}}) + (p_l \cdot p_{\bar{q}})(p_{\bar{l}} \cdot p_q)] \end{aligned} \quad (\text{A.28})$$

using the identity $g^{\mu\nu} g_{\mu\nu} = 4$.

From equations A.5 to A.8, $(p_l \cdot p_q) = (p_{\bar{l}} \cdot p_{\bar{q}}) = \frac{s'}{4}(1 + \cos \theta)$ and $(p_l \cdot p_{\bar{q}}) = (p_{\bar{l}} \cdot p_q) = \frac{s'}{4}(1 - \cos \theta)$. The cross section corresponding to the creation of a photon only is then (grouping equations A.12 and A.28):

$$\frac{d\sigma_\gamma}{d\Omega} = \frac{1}{3} \frac{e^4}{(4\pi)^2} Q_q^2 Q_l^2 \frac{1}{8s'} [(1 + \cos \theta)^2 + (1 - \cos \theta)^2] \quad (\text{A.29})$$

As previously said this cross section is indeed proportion to $\alpha^2 = (e^2/(4\pi))^2$. One also observes that this cross section remains the same for $\theta \rightarrow \pi - \theta$. The forward-backward asymmetry is then 0 for the photon contribution alone.

A.5.2 $|\mathcal{M}_Z|^2$

The calculation of $|\mathcal{M}_Z|^2$ is very similar to $|\mathcal{M}_\gamma|^2$. The only differences are the presence of the γ^5 matrix in the couplings and of a factor $\mathcal{R}^2$ coming from the Z propagator in front of the total amplitude. These differences have however a huge impact on the final result. The complex conjugate of $\mathcal{M}_Z$ is:

$$\mathcal{M}_Z^* = \frac{Q_l Q_q e^2}{s'} \mathcal{R}^* \bar{v}_{\bar{l}} \gamma^\nu (g_{V_l} + g_{A_l} \gamma^5) u_l \bar{u}_q \gamma_\nu (g_{V_q} + g_{A_q} \gamma^5) v_{\bar{q}} \quad (\text{A.30})$$

The amplitude squared contains again a part with the leptons and another with the quarks. The lepton part gives:

$$\sum_{s_l s_{\bar{l}}} \bar{u}_l \gamma^\mu (g_{V_l} + g_{A_l} \gamma^5) v_{\bar{l}} \bar{v}_l \gamma^\nu (g_{V_l} + g_{A_l} \gamma^5) u_l = \text{Tr} (p_{l_\rho} p_{\bar{l}_\sigma} \gamma^\rho \gamma^\mu (g_{V_l} + g_{A_l} \gamma^5) \gamma^\sigma \gamma^\nu (g_{V_l} + g_{A_l} \gamma^5)) \quad (\text{A.31})$$

From equation A.62 one sees that $\gamma^5 \gamma^\mu \gamma^\nu = \gamma^\mu \gamma^\nu \gamma^5$. We can therefore group the $(g_{V_l} + g_{A_l} \gamma^5)$ terms that will give $(g_{V_l} + g_{A_l} \gamma^5)^2 = ((g_{V_l}^2 + g_{A_l}^2) + 2g_{V_l} g_{A_l} \gamma^5)$ (using equation A.61). The same trick can be applied to the quarks brackets. Two traces must be multiplied:

$$\text{Tr} (p_{l_\rho} p_{\bar{l}_\sigma} \gamma^\rho \gamma^\mu \gamma^\sigma \gamma^\nu ((g_{V_l}^2 + g_{A_l}^2) + 2g_{V_l} g_{A_l} \gamma^5)) \cdot \text{Tr} (p_{q_\rho} p_{\bar{q}_\sigma} \gamma^\rho \gamma_\mu \gamma_\sigma \gamma_\nu ((g_{V_q}^2 + g_{A_q}^2) + 2g_{V_q} g_{A_q} \gamma^5)) \quad (\text{A.32})$$

This expression can be split into three terms:

- the term with no γ^5 (similar to the case of the photon, cf. equation A.28):

$$(g_{V_l}^2 + g_{A_l}^2)(g_{V_q}^2 + g_{A_q}^2)p_{l_\rho}p_{\bar{l}_\sigma}p_q^{\bar{\rho}}p_{\bar{q}}^{\bar{\sigma}}.Tr(\gamma^\rho\gamma^\mu\gamma^\sigma\gamma^\nu).Tr(\gamma_{\bar{\rho}}\gamma_\mu\gamma_{\bar{\sigma}}\gamma_\nu) \quad (A.33)$$

$$= (g_{V_l}^2 + g_{A_l}^2)(g_{V_q}^2 + g_{A_q}^2)16 \times 2 [(p_l \cdot p_q)(p_{\bar{l}}p_{\bar{q}}) + (p_l \cdot p_{\bar{q}})(p_{\bar{l}} \cdot p_q)] \quad (A.34)$$

- the term where γ^5 is present in only one of the traces. This trace is then proportional to $\epsilon^{\rho\mu\sigma\nu}$ which is antisymmetric under $\mu \rightarrow \nu$. The other is proportional to $Tr(\gamma_{\bar{\rho}}\gamma_\mu\gamma_{\bar{\sigma}}\gamma_\nu)$ which is symmetric under $\mu \rightarrow \nu$. The product of the two traces is therefore zero because of the sum over μ and ν .
- the term where γ^5 is present in both traces. The equation A.66 implies that this term is:

$$4(g_{V_l}g_{A_l}g_{V_q}g_{A_q})p_{l_\rho}p_{\bar{l}_\sigma}p_q^{\bar{\rho}}p_{\bar{l}}^{\bar{\sigma}}(-16)\epsilon^{\rho\mu\sigma\nu}\epsilon_{\bar{\rho}\mu\bar{\sigma}\nu} \quad (A.35)$$

$$= 4(g_{V_l}g_{A_l}g_{V_q}g_{A_q})p_{l_\rho}p_{\bar{l}_\sigma}p_q^{\bar{\rho}}p_{\bar{l}}^{\bar{\sigma}}(-16)\epsilon^{\rho\sigma\mu\nu}\epsilon_{\bar{\rho}\bar{\sigma}\mu\nu} \quad (A.36)$$

$$= 4(g_{V_l}g_{A_l}g_{V_q}g_{A_q})p_{l_\rho}p_{\bar{l}_\sigma}p_q^{\bar{\rho}}p_{\bar{l}}^{\bar{\sigma}}(-16)2(\delta_\sigma^\rho\delta_{\bar{\rho}}^\sigma - \delta_{\bar{\rho}}^\rho\delta_\sigma^\sigma) \quad (A.37)$$

$$= 128(g_{V_l}g_{A_l}g_{V_q}g_{A_q})[(p_l \cdot p_q)(p_{\bar{l}}p_{\bar{q}}) - (p_l \cdot p_{\bar{q}})(p_{\bar{l}} \cdot p_q)] \quad (A.38)$$

where we used the equation A.67 in the third line.

We can now compute the cross section corresponding to the creation of a Z boson only (grouping equations A.12, A.34 and A.38):

$$\frac{d\sigma_Z}{d\Omega} = \frac{1}{3} \frac{e^4}{(4\pi)^2} Q_q^2 Q_l^2 \frac{1}{8s'} |\mathcal{R}|^2 [c_{1,Z}(1 + \cos\theta)^2 + c_{2,Z}(1 - \cos\theta)^2] \quad (A.39)$$

where

$$c_{1,Z} = ((g_{V_l}^2 + g_{A_l}^2)(g_{V_q}^2 + g_{A_q}^2) + 4g_{V_l}g_{A_l}g_{V_q}g_{A_q}) \quad (A.40)$$

$$c_{2,Z} = ((g_{V_l}^2 + g_{A_l}^2)(g_{V_q}^2 + g_{A_q}^2) - 4g_{V_l}g_{A_l}g_{V_q}g_{A_q}) \quad (A.41)$$

Because $c_{1,Z}$ and $c_{2,Z}$ are different, the cross section is not invariant under $\theta \rightarrow \pi - \theta$. The Z boson alone can therefore introduce a forward-backward asymmetry under the condition $g_{V_l}g_{A_l}g_{V_q}g_{A_q} \neq 0$, that is, no coupling vanishes. It is interesting to note that this term does not depend on the center of mass energy.

A.5.3 $\mathcal{M}_\gamma \mathcal{M}_Z^* + c.c$

The computation for the interference follows the same procedure than for the two other terms. The only difference is that this time, we mix the Z boson and photon couplings and that the final product will be proportional to $\text{Re}(\mathcal{R})$.

Let us start with $\mathcal{M}_\gamma \mathcal{M}_Z^*$. The lepton part of the matrix element gives, applying the same rules than previously:

$$\sum_{s_l s_{\bar{l}}} \bar{u}_l \gamma^\mu v_{\bar{l}} \bar{v}_l \gamma^\nu (g_{V_l} + g_{A_l} \gamma^5) u_l = p_{l_\rho} p_{\bar{l}_\sigma} . Tr(\gamma^\rho \gamma^\mu (g_{V_l} + g_{A_l} \gamma^5) \gamma^\sigma \gamma^\nu) \quad (A.42)$$

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This must be multiplied by the trace coming from the quark part. The reasoning is identical to the case of the $|\mathcal{M}_Z|^2$. One finds:

$$\frac{1}{4} \sum_{s_l s_{\bar{l}} s_q s_{\bar{q}}} \mathcal{M}_\gamma \mathcal{M}_Z^* = \frac{e^4}{4} Q_q^2 Q_l^2 \mathcal{R} [c_{1,int}/2(1 + \cos \theta)^2 + c_{2,int}/2(1 - \cos \theta)^2] \quad (\text{A.43})$$

where

$$c_{1,int} = 2(g_{V_l} g_{V_q} + g_{A_l} g_{A_q}) \quad (\text{A.44})$$

$$c_{2,int} = 2(g_{V_l} g_{V_q} - g_{A_l} g_{A_q}) \quad (\text{A.45})$$

The complex conjugate term is the same with $\mathcal{R} \rightarrow \mathcal{R}^*$ and the cross section from the interference is:

$$\frac{d\sigma_{int}}{d\Omega} = \frac{1}{3} \frac{e^4}{(4\pi)^2} Q_q^2 Q_l^2 \frac{1}{8s'} \text{Re}(\mathcal{R}) [c_{1,int}(1 + \cos \theta)^2 + c_{2,int}(1 - \cos \theta)^2] \quad (\text{A.46})$$

For the interference, the condition for a forward-backward asymmetry to arise is only that g_{A_l} and g_{A_q} are non zero. This is indeed expected. We see from the computation $|\mathcal{M}_Z|^2$ that what produces the asymmetry is the mixture of an axial and a vectorial coupling. In the interference, the vectorial component is already obtained from the photon. Similarly to the Z boson case, the asymmetry from this single contribution does not depend of $\sqrt{s'}$.

A.6 The final cross section

The total angular differential cross section is the sum of equations A.29, A.39 and A.46. To make the asymmetry more explicit we have re-arranged the angular terms:

$$\frac{d\sigma_{\gamma/Z}}{d\Omega} = \frac{1}{3} \frac{e^4}{(4\pi)^2} Q_q^2 Q_l^2 \frac{1}{4s'} [c_1(1 + \cos^2 \theta) + c_2 \cos \theta] \quad (\text{A.47})$$

where:

$$c_1 = 1 + 2\text{Re}(\mathcal{R})g_{V_l}g_{V_q} + |\mathcal{R}|^2(g_{V_l}^2 + g_{A_l}^2)(g_{V_q}^2 + g_{A_q}^2) \quad (\text{A.48})$$

$$c_2 = 4\text{Re}(\mathcal{R})g_{A_l}g_{A_q} + 8|\mathcal{R}|^2g_{V_l}g_{A_l}g_{V_q}g_{A_q} \quad (\text{A.49})$$

We can now compute the total cross section:

$$\sigma = \int_{\Omega} \frac{d\sigma_{\gamma+Z}}{d\Omega} d\Omega = \int_0^{2\pi} d\phi \int_0^\pi \sin\theta d\theta \frac{\alpha^2 Q_q^2 Q_l^2}{12s'} [c_1(1 + \cos^2 \theta) + c_2 \cos \theta] = \frac{4\pi}{9} \frac{\alpha^2 Q_q^2}{s'} \quad (\text{A.50})$$

where we used again $\alpha = e^2/(4\pi)$. The forward and backward cross sections are respectively given by:

$$\sigma_F = \int_0^{2\pi} d\phi \int_0^{\pi/2} \sin\theta d\theta \frac{\alpha^2 Q_q^2 Q_l^2}{4s'} [c_1(1 + \cos^2 \theta) + c_2 \cos \theta] = \alpha^2 Q_q^2 Q_l^2 \frac{\pi}{6s'} \left(\frac{4}{3}c_1 + \frac{1}{2}c_2 \right) \quad (\text{A.51})$$

$$\sigma_B = \int_0^{2\pi} d\phi \int_{\pi/2}^{\pi} \sin\theta d\theta \frac{\alpha^2 Q_q^2 Q_l^2}{4s'} [c_1(1 + \cos^2 \theta) + c_2 \cos \theta] = \alpha^2 Q_q^2 Q_l^2 \frac{\pi}{6s'} \left(\frac{4}{3} c_1 - \frac{1}{2} c_2 \right) \quad (\text{A.52})$$

And finally the forward-backward asymmetry is:

$$A_{FB} = \frac{\sigma_F - \sigma_B}{\sigma_F + \sigma_B} = \frac{3}{8} \frac{c_2}{c_1} \quad (\text{A.53})$$

The asymmetry now depends of $\sqrt{s'}$ because of the presence of $\mathcal{R}$ in c_1 and c_2 . It vanishes if $c_2 = 0$, that is:

$$\text{Re}(\mathcal{R}) + 2|\mathcal{R}|^2 g_{V_l} g_{V_q} = 0 \quad (\text{A.54})$$

One has, from equation (A.17):

$$\text{Re}(\mathcal{R}) = \text{Re} \left(\frac{1}{Q_l Q_q \sin^2 2\theta_W} \frac{1}{1 - M_Z^2/s' + i\Gamma_Z/M_Z} \right) \quad (\text{A.55})$$

$$= \frac{1}{Q_l Q_q \sin^2 2\theta_W} \frac{1 - M_Z^2/s'}{(1 - M_Z^2/s')^2 + (\Gamma_Z/M_Z)^2} \quad (\text{A.56})$$

$$|\mathcal{R}|^2 = \left(\frac{1}{Q_l Q_q \sin^2 2\theta_W} \right)^2 \frac{1}{(1 - M_Z^2/s')^2 + (\Gamma_Z/M_Z)^2} \quad (\text{A.57})$$

and thus, (A.54) gives:

$$(1 - M_Z^2/s') + \frac{2g_{V_l} g_{V_q}}{Q_l Q_q \sin^2 2\theta_W} = 0 \quad (\text{A.58})$$

The asymmetry vanishes for $s' = M_Z^2/(1 + 2g_{V_l} g_{V_q}/(Q_l Q_q \sin^2 2\theta_W))$. Numerically, $\sin^2 2\theta_W \approx 0.75$, $g_{V_l}/Q_l \approx 0.02$, $g_{V_u}/Q_u \approx 0.3$ and $g_{V_d}/Q_d \approx 1$. The term $2g_{V_l} g_{V_q}/(Q_l Q_q \sin^2 2\theta_W)$ is then always positive and $\ll 1$. As a consequence the energy where the asymmetry vanishes is closer (but a bit smaller) than the Z peak mass. As $g_{V_d}/Q_d > g_{V_u}/Q_u$, the 0 asymmetry energy is slightly smaller for down quarks than for up quarks.

A.7 The Feynman rules

The Feynman rules (in the Feynman gauge) used in this appendix are shown in table A.1. The value of the charge and the third component of the weak isospin are listed in Table A.2.

A.8 Properties of the γ matrices and the spinors

The γ^μ matrices are a set of 4 x 4 matrices generating the Dirac algebra and are therefore defined by the anticommutation rules:

$$\{\gamma^\mu, \gamma^\nu\} \equiv 2g^{\mu\nu} . I_{4 \times 4} \quad (\text{A.59})$$

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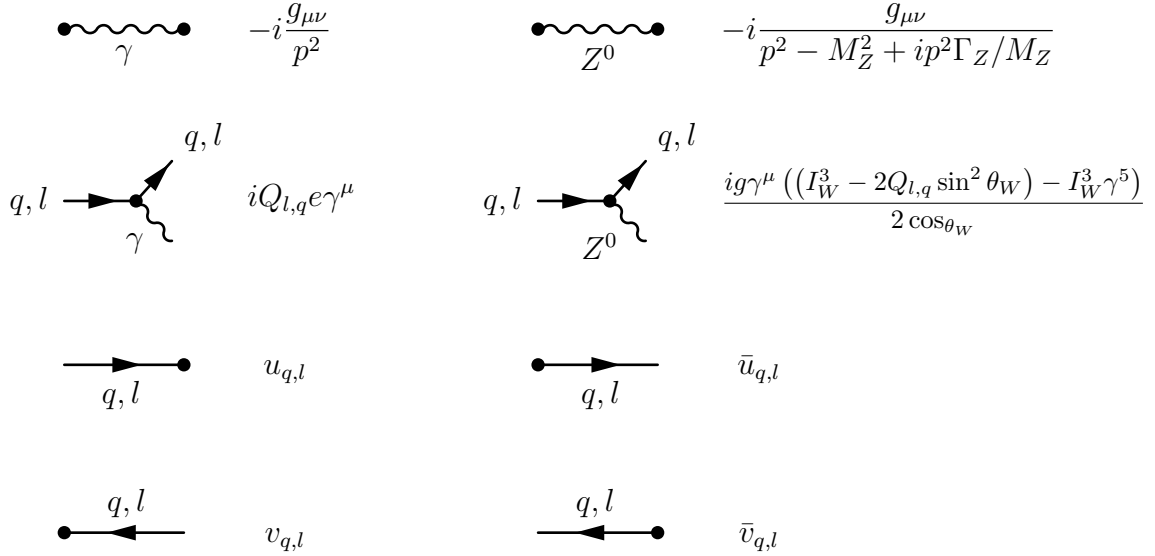


Table A.1: The Feynman rules used in the calculation of the Drell-Yan cross section.

The γ^5 matrix is defined as $\gamma^5 \equiv i\gamma^0\gamma^1\gamma^2\gamma^3$. We will use the chiral basis in which the γ matrices have the following properties:

$$(\gamma^0)^2 = I_{4 \times 4}; \quad (\text{A.60})$$

$$(\gamma^5)^2 = I_{4 \times 4}; \quad (\text{A.61})$$

$$\{\gamma^5, \gamma^\mu\} = 0; \quad (\text{A.62})$$

$$(\gamma^\mu)^\dagger = \gamma^0 \gamma^\mu \gamma^0; \quad (\text{A.63})$$

$$(\gamma^5)^\dagger = \gamma^5; \quad (\text{A.64})$$

$$\text{Tr}(\gamma^\alpha \gamma^\beta \gamma^\mu \gamma^\nu) = 4(g^{\alpha\beta} g^{\mu\nu} - g^{\alpha\mu} g^{\beta\nu} + g^{\alpha\nu} g^{\beta\mu}); \quad (\text{A.65})$$

$$\text{Tr}(\gamma^\alpha \gamma^\beta \gamma^\mu \gamma^\nu \gamma^5) = -4i\epsilon^{\alpha\beta\mu\nu}. \quad (\text{A.66})$$

where ϵ is the totally antisymmetric tensor and is defined such that $\epsilon^{0123} = 1$. It has the property:

$$\epsilon^{\alpha\beta\mu\nu} \epsilon_{\bar{\alpha}\bar{\beta}\bar{\mu}\bar{\nu}} = 2(\delta_{\bar{\beta}}^\alpha \delta_{\bar{\alpha}}^\beta - \delta_{\bar{\alpha}}^\alpha \delta_{\bar{\beta}}^\beta) \quad (\text{A.67})$$

The u and v spinors involved verify the following identities:

$$\sum_s u_\alpha(p, s) \bar{u}_\beta(p, s) = (p_\rho \gamma^\rho + m I_{4 \times 4})_{\alpha\beta}; \quad (\text{A.68})$$

$$\sum_s v_\alpha(p, s) \bar{v}_\beta(p, s) = (p_\rho \gamma^\rho - m I_{4 \times 4})_{\alpha\beta}. \quad (\text{A.69})$$

	Q	I_W^3
ν_e, ν_μ, ν_τ	0	$\frac{1}{2}$
e, μ, τ	-1	$-\frac{1}{2}$
u, c, t	$\frac{2}{3}$	$\frac{1}{2}$
d, s, b	$-\frac{1}{3}$	$-\frac{1}{2}$

Table A.2: Charge (Q) and third component of the weak isospin (I_W^3) for the fermions in the Standard Model.

A. DRELL-YAN CROSS SECTION CALCULATION

Appendix B

Performance of the various conditions entering the high E_T electron selection

This appendix presents the performance in terms of electron and fake efficiencies of the conditions applied on the variables entering the high E_T electron selection (*cf.* table 4.4). Each of the conditions was initially tuned such that it keeps around 99% of the simulated electrons reconstructed as GSF electron candidates with $E_T > 100$ GeV.

The discrimination power of each variable is quantified in section B.1 where the electron efficiency is plotted as a function of the fake efficiency for various cut values. The resulting curves are shown for electrons with $E_T > 35$ GeV and $E_T > 100$ GeV separately.

The efficiency of each of the conditions is then drawn as a function of E_T for GSF electron candidates (“single efficiency”) and for GSF electron candidates having already passed the other conditions of the selection (“(N-1) efficiency”) in section B.2. The two sets of plots are complementary: the single efficiencies indicate the intrinsic performance of each condition whereas the (N-1) efficiencies show that the correlation between the various conditions is limited.

In this appendix, the same samples and event selection than in section 4.3 are used. The electron efficiency is calculated from GSF electron candidates found in a Drell-Yan simulated sample and matching a generated electron. Since the selection mainly aims at rejecting jets, fakes are defined here as GSF electron candidates from a QCD sample. For two variables however (d_{xy} and the number of missing hits in the innermost layers of the tracker), results for a γ +jets sample are also shown in section B.1 since those two variables are used to reject photons that converted in the tracker material rather than jets.

B. PERFORMANCE OF THE VARIOUS CONDITIONS ENTERING THE HIGH E_T ELECTRON SELECTION

B.1 Electron versus background efficiency curves

Figures B.1, B.2 and B.3 (B.4, B.5 and B.6) present, for GSF electron candidates with $E_T > 35$ GeV ($E_T > 100$ GeV), the electron efficiency versus jet efficiency curve for each of the variables forming the high E_T electron selection. For the $|\Delta\eta_{in}|$, $|\Delta\phi_{in}|$, H/E , $\sigma_{in\eta}$ (endcaps), tracker isolation and $|d_{xy}|$ variables, the construction is well defined: various cut values are scanned and the corresponding efficiencies are computed. The transverse shower shape condition in the barrel consists of two alternative cuts and the construction of the curve is not unique. The curves in figures B.2 and B.5 (top left) show the efficiency obtained when $E_{1\times5}/E_{1\times5}$ is requested to be greater than 0.83 while the $E_{1\times5}/E_{1\times5}$ cut value varies. For the calorimeter isolation, the chosen floating parameter is the cut value on $2 + 0.03E_T - 0.28\rho$ (barrel) or on $Max(1 + 0.03E_T - 0.28\rho, 2.5 - 0.28\rho)$ (endcaps).

In figures B.3 and B.6, the curve for the $|d_{xy}|$ variable is also presented for a γ +jets sample. The figures also show the efficiencies for a missing hit cut value of 0 (leftmost dot) or ≤ 1 (rightmost dot).

The red dot in each figure indicates the values corresponding to the cut used in the high E_T electron selection.

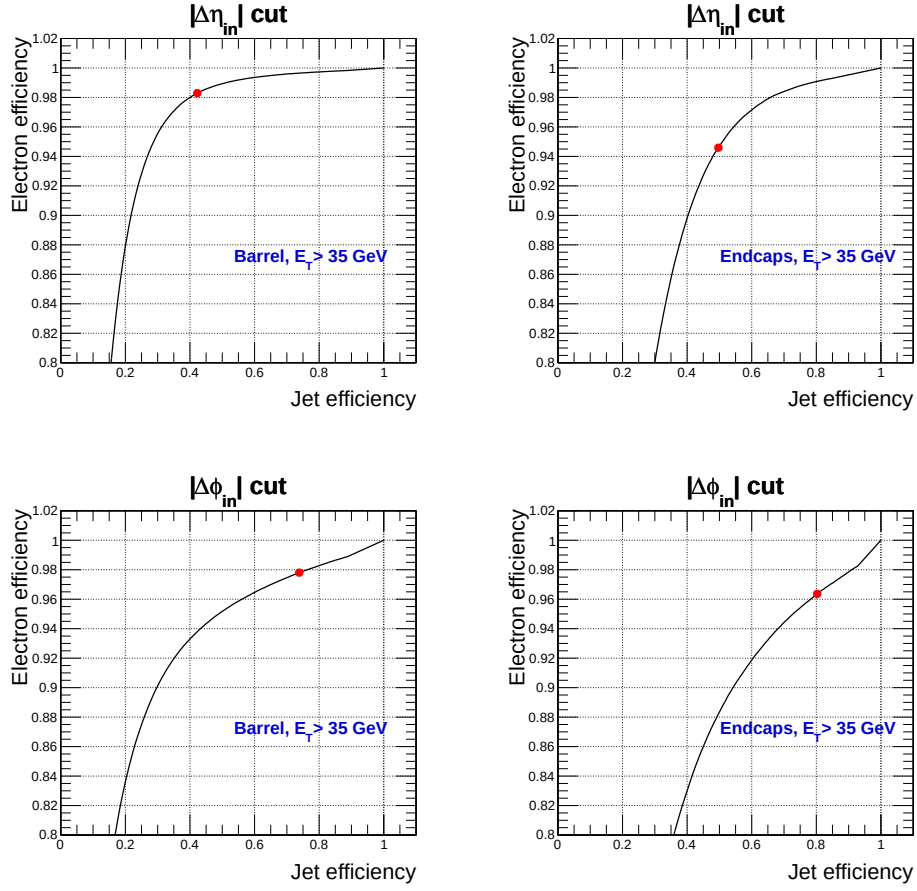


Figure B.1: Electron versus jet efficiency for GSF electron candidates with $E_T > 35$ GeV when applying a cut on the $|\Delta\eta_{in}|$ (top) or $|\Delta\phi_{in}|$ (bottom) variable. Results are shown separately for electrons reconstructed in the barrel (left) or in the endcaps (right). The red dot indicates the values corresponding to the cut used in the high E_T electron selection.

B. PERFORMANCE OF THE VARIOUS CONDITIONS ENTERING THE HIGH E_T ELECTRON SELECTION

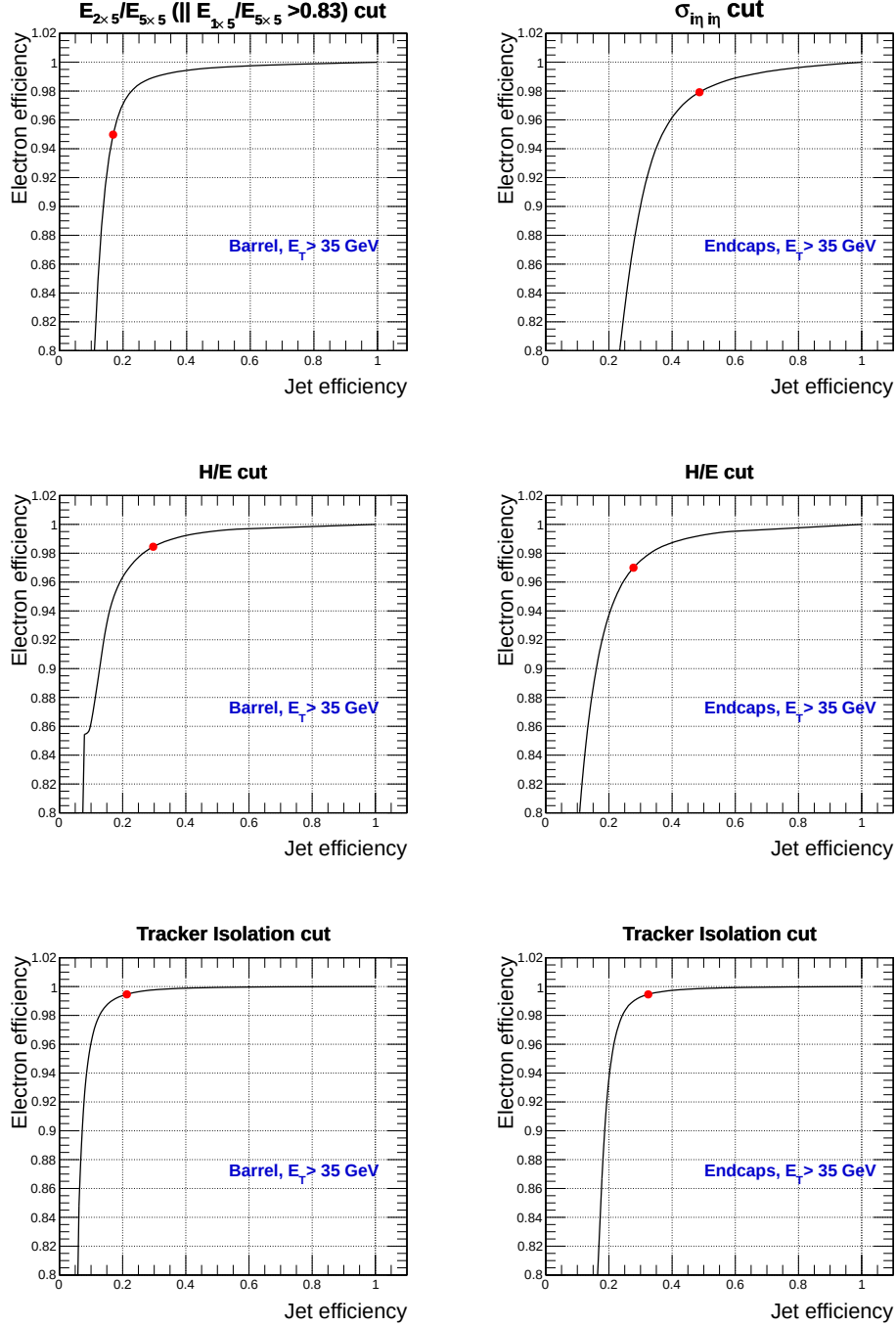


Figure B.2: Electron versus jet efficiency for GSF electron candidates with $E_T > 35$ GeV when applying a cut on the transverse shower shape (top), on H/E (middle) or on the tracker isolation variable (bottom). Results are shown separately for electrons reconstructed in the barrel (left) or in the endcaps (right). The red dot indicates the values corresponding to the cut used in the high E_T electron selection.

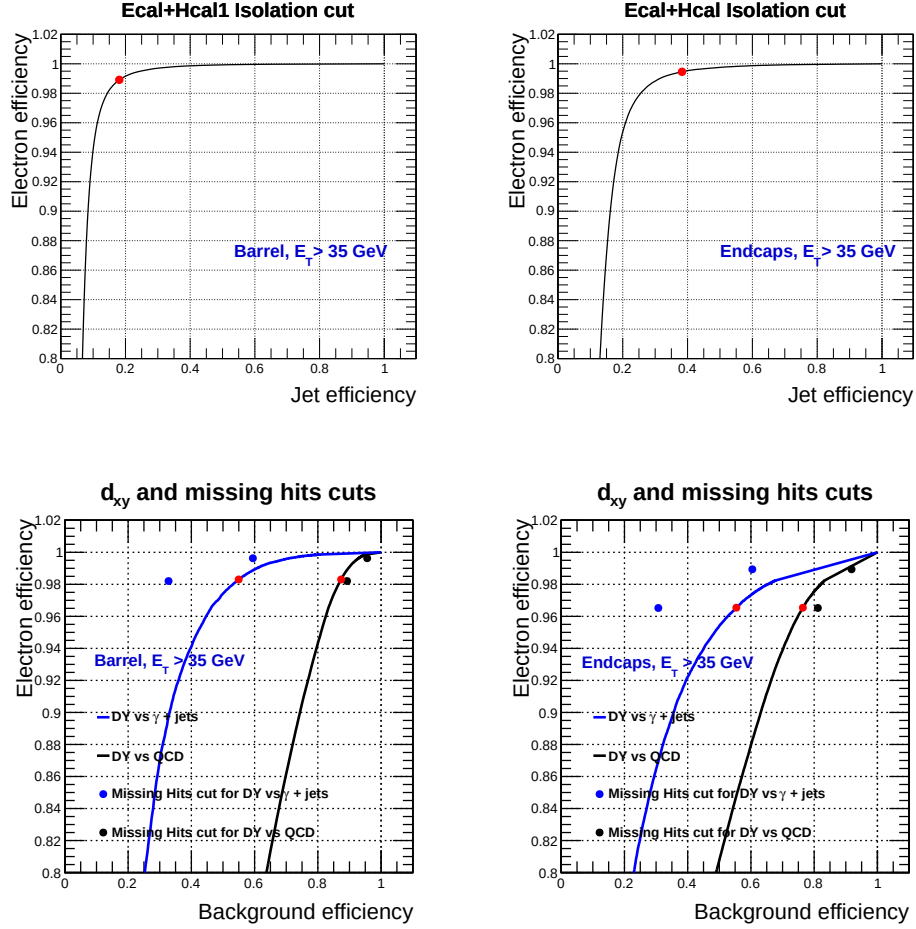


Figure B.3: Electron versus jet efficiency for GSF electron candidates with $E_T > 35$ GeV when applying a cut on the ECAL+HCAL1 isolation (top) or d_{xy} variable (bottom). The bottom plots also indicate the efficiency values corresponding to a cut on the number of missing hits equal to 0 (leftmost dot) or smaller or equal to 1 (rightmost dot). In these plots, the corresponding curve and dots for GSF electron candidates coming from a γ +jets sample instead of a QCD sample are shown in blue. Results are shown separately for electrons reconstructed in the barrel (left) or in the endcaps (right). The red dot indicates the values corresponding to the cut used in the high E_T electron selection.

B. PERFORMANCE OF THE VARIOUS CONDITIONS ENTERING THE HIGH E_T ELECTRON SELECTION

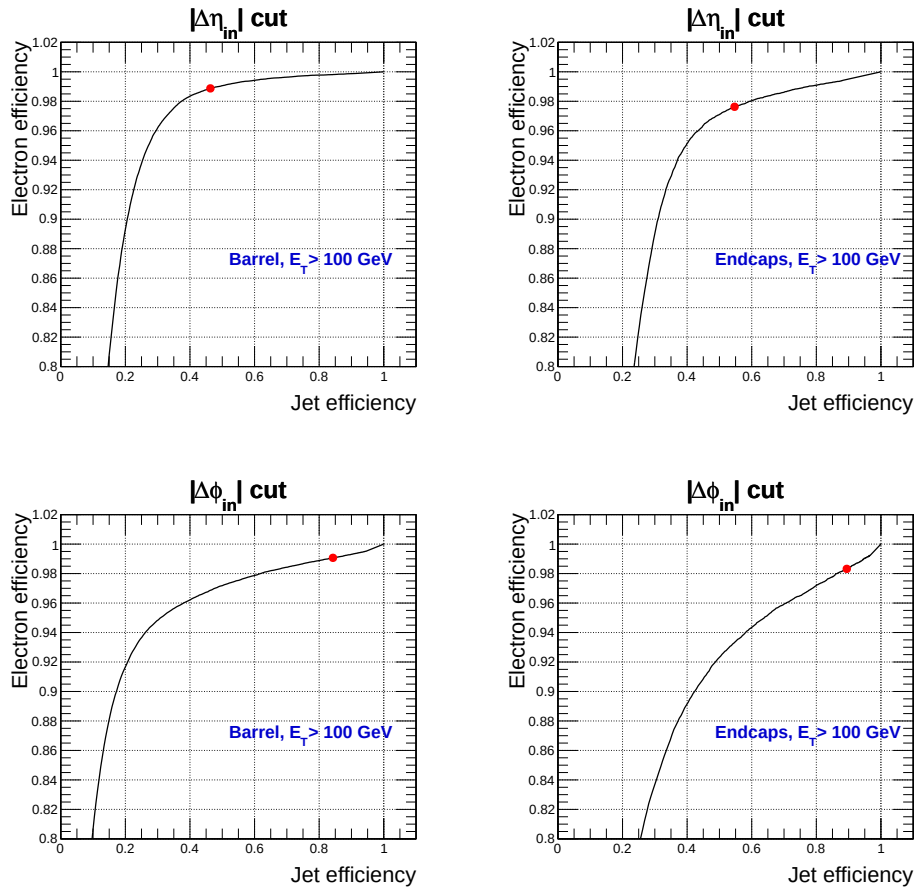


Figure B.4: Electron versus jet efficiency for GSF electron candidates with $E_T > 100$ GeV when applying a cut on the $|\Delta\eta_{in}|$ (top) or $|\Delta\phi_{in}|$ (bottom) variable. Results are shown separately for electrons reconstructed in the barrel (left) or in the endcaps (right). The red dot indicates the values corresponding to the cut used in the high E_T electron selection.

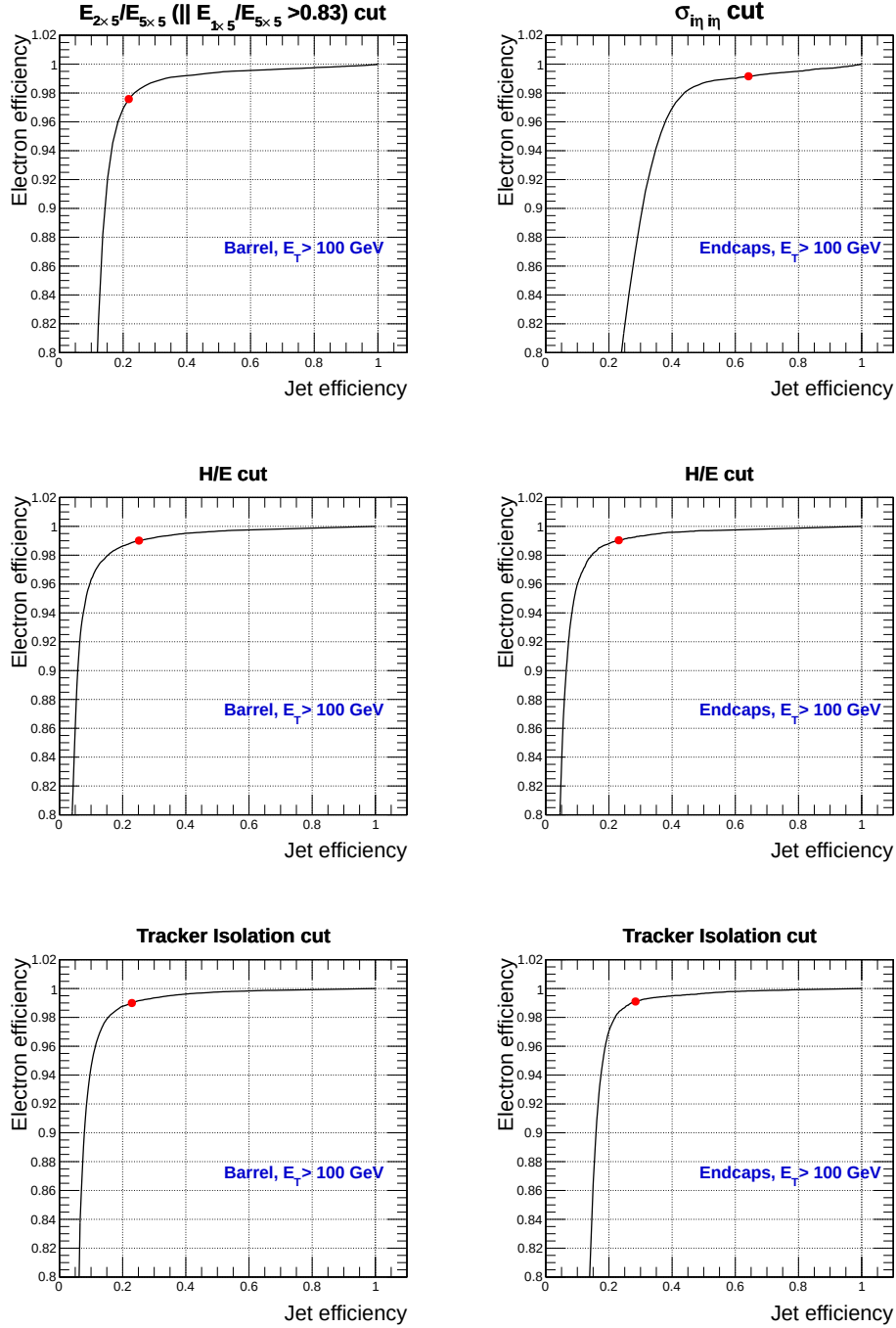


Figure B.5: Electron versus jet efficiency for GSF electron candidates with $E_T > 100$ GeV when applying a cut on the transverse shower shape (top), on H/E (middle) or on the tracker isolation variable (bottom). Results are shown separately for electrons reconstructed in the barrel (left) or in the endcaps (right). The red dot indicates the values corresponding to the cut used in the high E_T electron selection.

B. PERFORMANCE OF THE VARIOUS CONDITIONS ENTERING THE HIGH E_T ELECTRON SELECTION

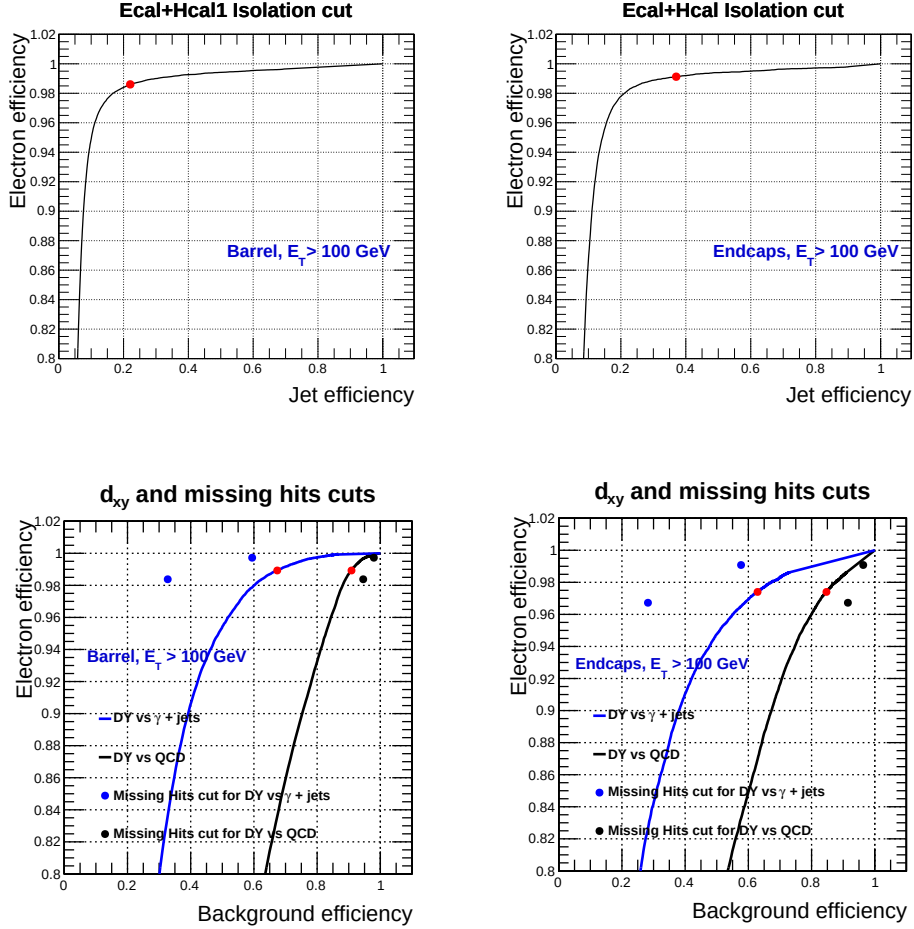


Figure B.6: Electron versus jet efficiency for GSF electron candidates with $E_T > 100$ GeV when applying a cut on the ECAL+HCAL1 isolation (top) or d_{xy} variable (bottom). The bottom plots also indicate the efficiency values corresponding to a cut on the number of missing hits equal to 0 (leftmost dot) or smaller or equal to 1 (rightmost dot). In these plots, the corresponding curve and dots for GSF electron candidates coming from a γ +jets sample instead of a QCD sample are shown in blue. Results are shown separately for electrons reconstructed in the barrel (left) or in the endcaps (right). The red dot indicates the values corresponding to the cut used in the high E_T electron selection.

B.2 Efficiencies of the various cuts as a function of E_T

B.2.1 Single efficiencies

Figures B.7, B.8 and B.9 present the efficiency of each of the various cuts entering the high E_T electron selection as a function of E_T for GSF electron candidates, from a Drell-Yan simulated sample (blue) and from a QCD simulated sample (orange) separately for barrel and endcaps electrons.

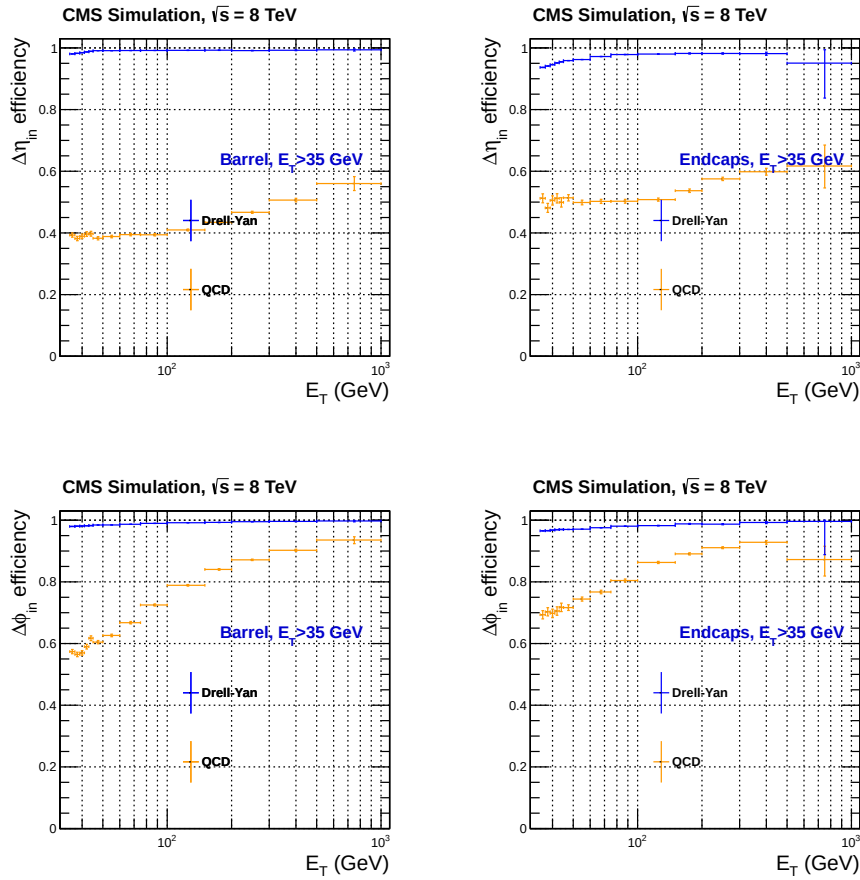


Figure B.7: Efficiency of the $|\Delta\eta_{in}|$ (top) and $|\Delta\phi_{in}|$ (bottom) cuts used in the high E_T electron selection as a function of E_T for GSF electron candidates, from a Drell-Yan simulated sample (blue) and from a QCD simulated sample (orange) separately for barrel and endcaps electrons.

B. PERFORMANCE OF THE VARIOUS CONDITIONS ENTERING THE HIGH E_T ELECTRON SELECTION

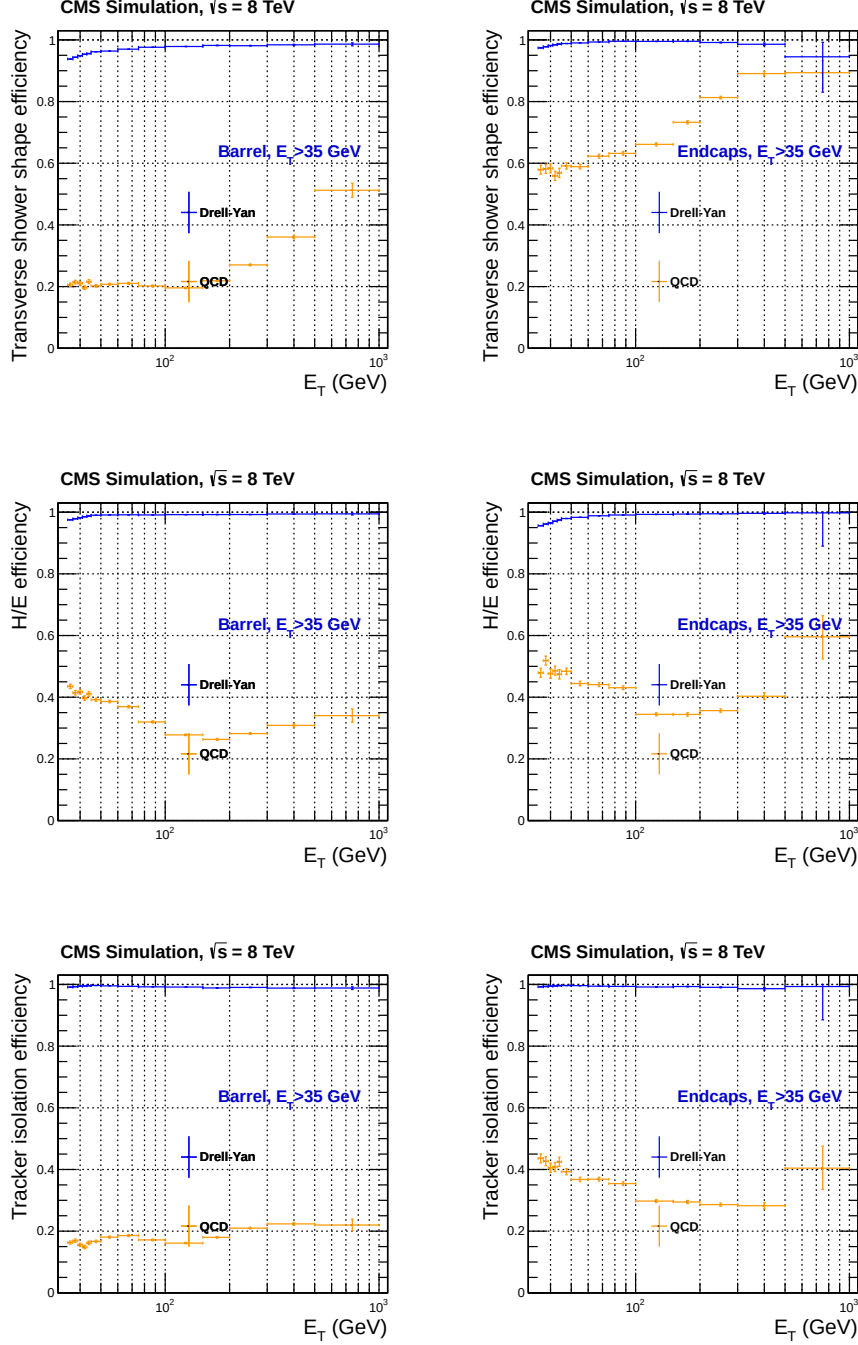


Figure B.8: Efficiency of the transverse shower shape (top), H/E (middle) and the tracker isolation (bottom) cuts used in the high E_T electron selection as a function of E_T for GSF electron candidates, from a Drell-Yan simulated sample (blue) and from a QCD simulated sample (orange) separately for barrel and endcaps electrons.

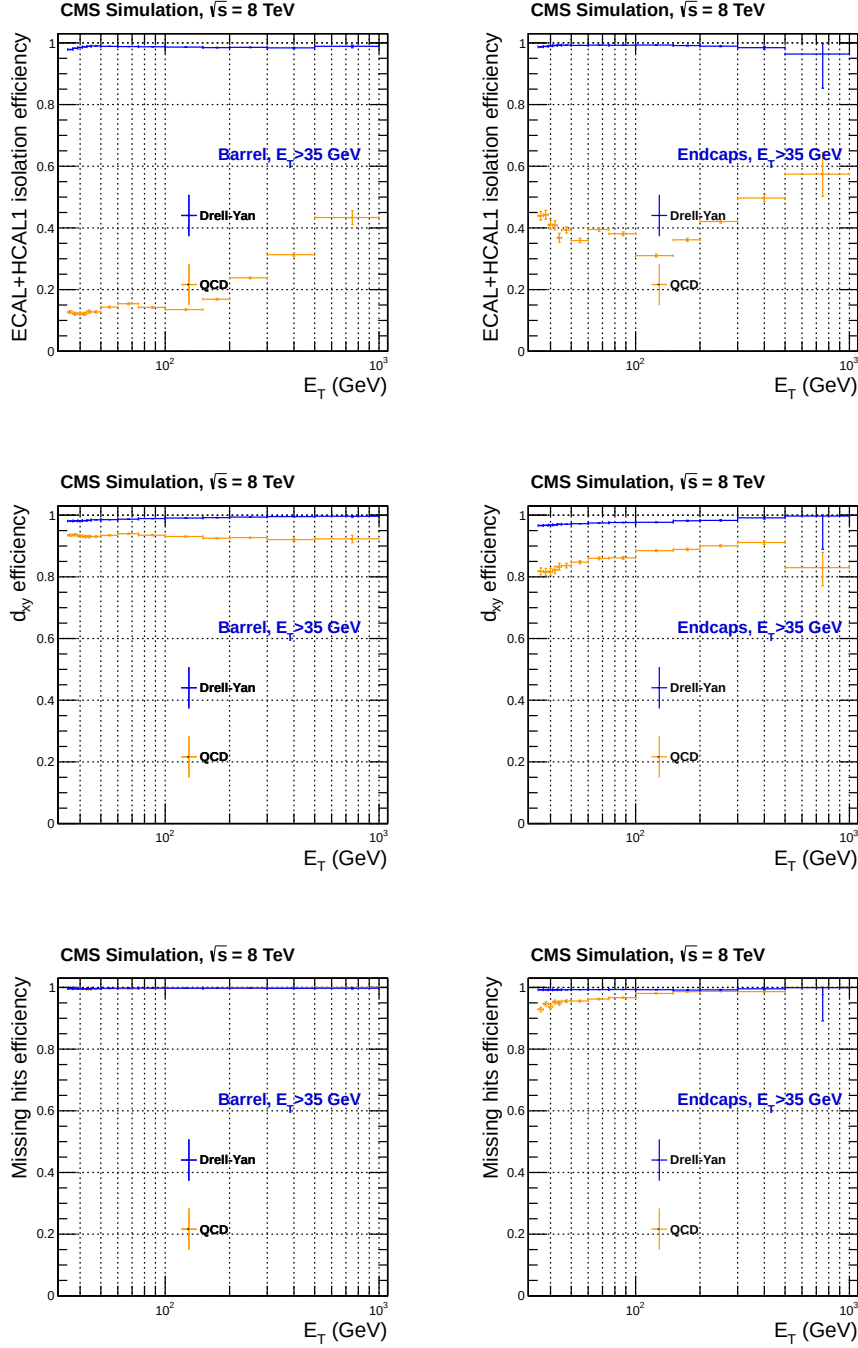


Figure B.9: Efficiency of the ECAL+HCAL1 isolation (top), $|d_{xy}|$ and the number of missing hits (bottom) cuts used in the high E_T electron selection as a function of E_T for GSF electron candidates, from a Drell-Yan simulated sample (blue) and from a QCD simulated sample (orange) separately for barrel and endcaps electrons.

B. PERFORMANCE OF THE VARIOUS CONDITIONS ENTERING THE HIGH E_T ELECTRON SELECTION

B.2.2 (N-1) efficiencies

Figures B.10, B.11 and B.12 present the efficiency of each of the various cuts entering the high E_T electron selection as a function of E_T for GSF electron candidates passing all the other cuts of the selection, from a Drell-Yan simulated sample (blue) and from a QCD simulated sample (orange) separately for barrel and endcaps electrons.

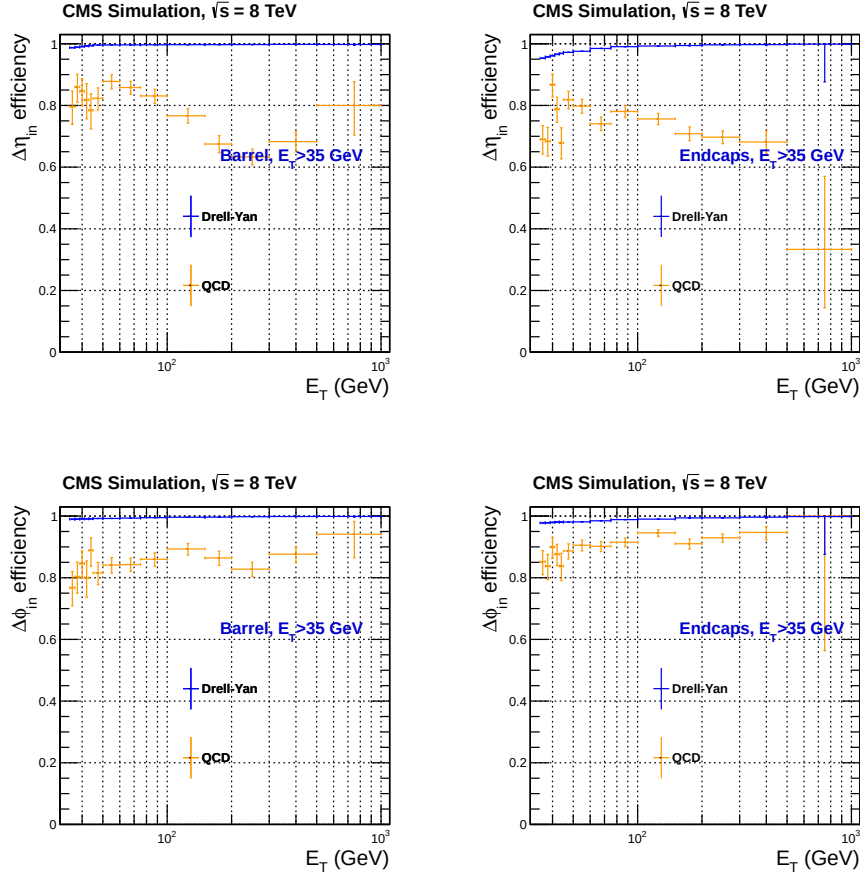


Figure B.10: Efficiency of the $|\Delta\eta_{in}|$ (top) and $|\Delta\phi_{in}|$ (bottom) cuts used in the high E_T electron selection as a function of E_T for GSF electron candidates passing all the other cuts of the selection, from a Drell-Yan simulated sample (blue) and from a QCD simulated sample (orange) separately for barrel and endcaps electrons.

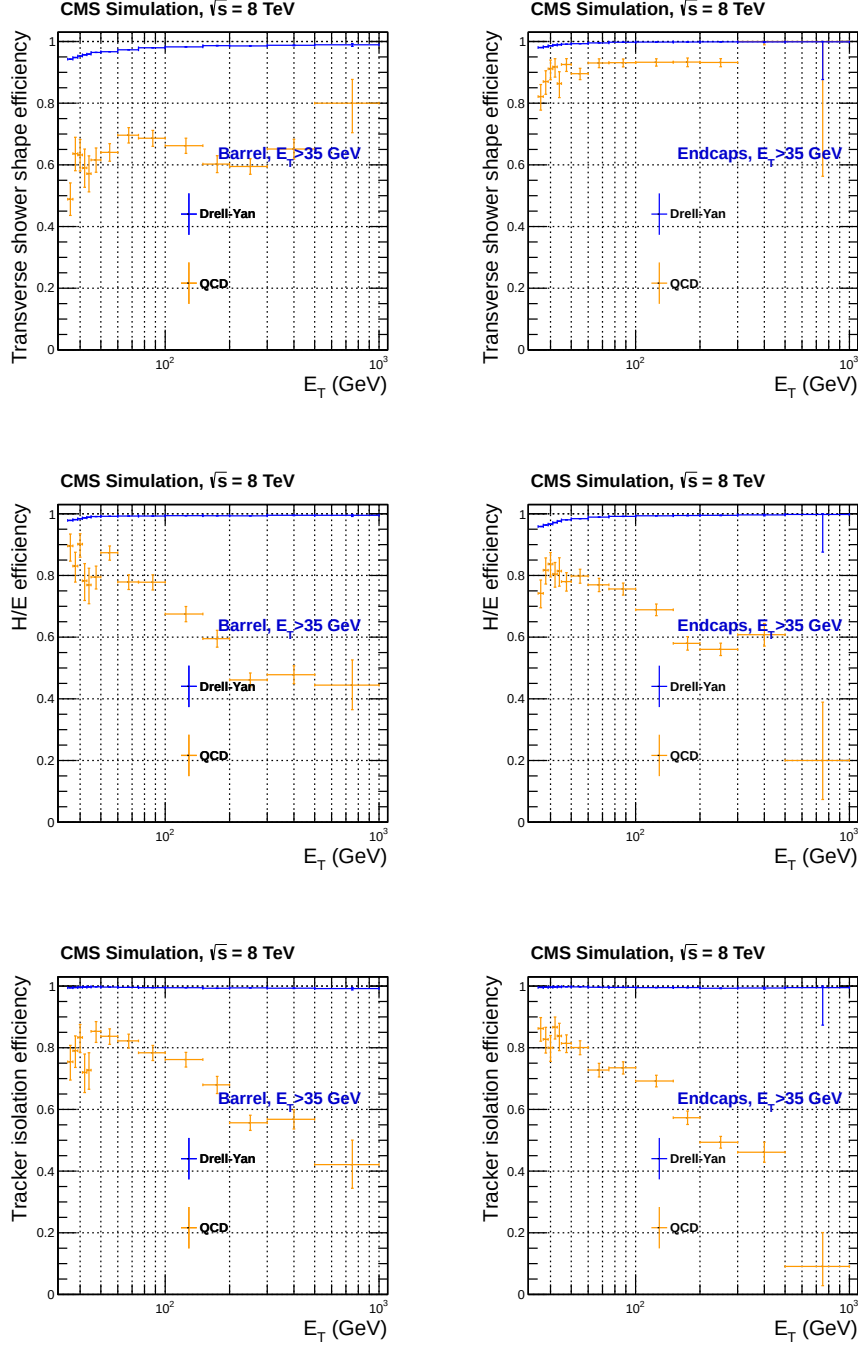


Figure B.11: Efficiency of the transverse shower shape (top), H/E (middle) and the tracker isolation (bottom) cuts used in the high E_T electron selection as a function of E_T for GSF electron candidates passing all the other cuts of the selection, from a Drell-Yan simulated sample (blue) and from a QCD simulated sample (orange) separately for barrel and endcaps electrons.

B. PERFORMANCE OF THE VARIOUS CONDITIONS ENTERING THE HIGH E_T ELECTRON SELECTION

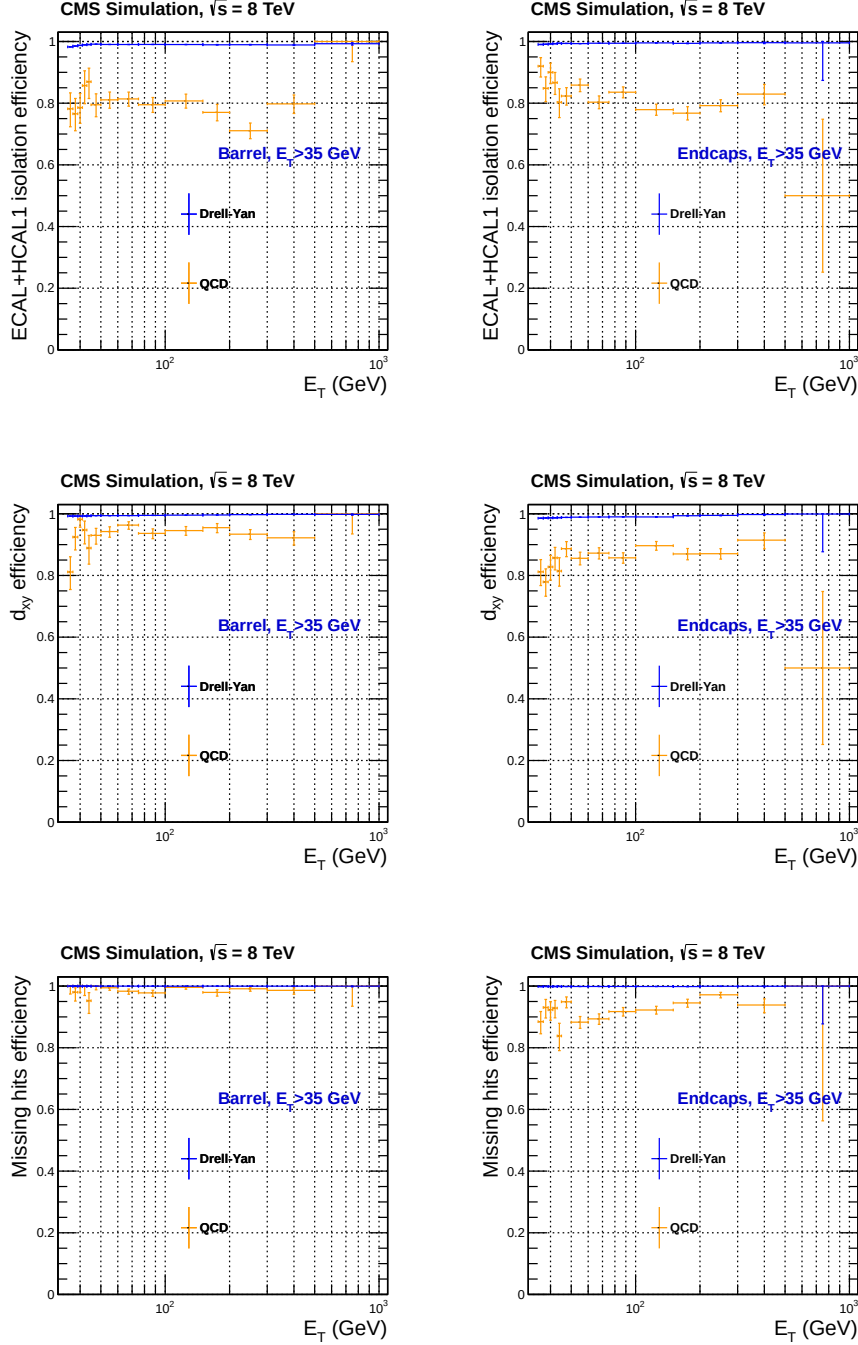


Figure B.12: Efficiency of the ECAL+HCAL1 isolation (top), $|d_{xy}|$ and the number of missing hits (bottom) cuts used in the high E_T electron selection as a function of E_T for GSF electron candidates passing all the other cuts of the selection, from a Drell-Yan simulated sample (blue) and from a QCD simulated sample (orange) separately for barrel and endcaps electrons.

Appendix C

Effective areas measurement at $\sqrt{s} = 7$ TeV

In chapter 4 the reconstruction and identification of high energy electrons with the CMS detector is presented. The starting point is the association of an energy deposit in the electromagnetic calorimeter with a track in the tracker. In order to disentangle electrons from other particles such as jets or photons, a series of cuts on various variables describing the shower shape, the isolation,... are then applied.

As explained in section 4.4, the distributions of some of those variables are spoiled by the presence of energy deposits arising from pile-up events. A technique using the measured average transverse energy density coming from pile-up collisions in the event is therefore used to correct the variables studied. The results presented in chapter 4 were obtained from simulated samples generated at $\sqrt{s} = 8$ TeV and are quite recent.

The correction factors (effective areas) were initially calculated using simulations generated at $\sqrt{s} = 7$ TeV and it is those corrections that were applied for the analysis of the 2012 data. This appendix presents those original results.

C.1 Method

The method is the same as described in section 4.4.2. The contribution from the pile-up to any energy measurement V_{raw} in the calorimeters can be estimated by multiplying the average transverse energy density coming from the pile-up in the event, denoted ρ , by an effective area (E_A). This effective area can depend on η and is estimated as the ratio of the linear evolution of value of V_{raw} below which 95% of the events are found over the linear evolution ρ as a function of the number of reconstructed vertices in the event. The corrected variable V_{corr} is then expressed as:

$$V_{corr} = V_{raw} - \rho \times E_A \quad (C.1)$$

C.2 ECAL+HCAL1 isolation correction

Figure C.1 shows the ECAL+HCAL1 isolation variable (*cf.* section 4.3.7) for barrel and endcaps GSF electron candidates from simulated Drell-Yan events, as a function of the

C. EFFECTIVE AREAS MEASUREMENT AT $\sqrt{s} = 7$ TEV

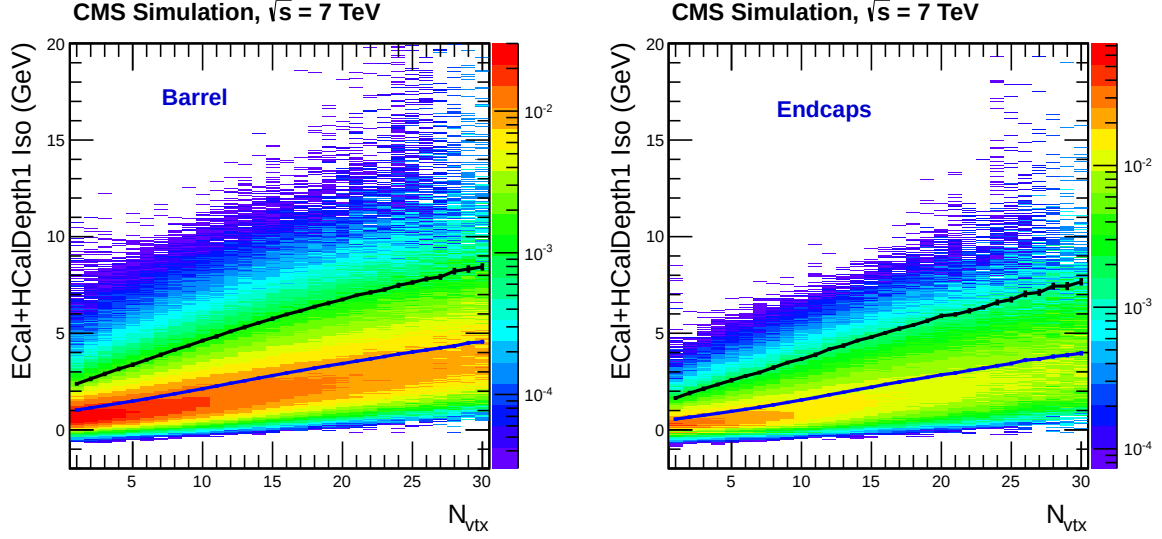


Figure C.1: The distribution of the ECAL+HCAL1 Isolation versus the number of primary vertices in the event for electrons with $E_T > 35$ GeV coming from a Drell-Yan simulated sample at $\sqrt{s} = 7$ TeV. Each slice in constant number of vertices is normalized to 1. For a given number of vertexes, the blue dots indicate the mean value of the projected distribution while the black dots indicate the value below which 95% of the events are contained.

number of vertices in the event. The effective area for the ECAL and HCAL1 isolation is computed separately and together for the two calorimeters and is shown in figure C.2 as a function of the supercluster pseudorapidity. Since the sum of the two isolation is roughly constant versus η , a constant effective area of 0.28 is used in the correction.

C.3 H/E correction

For this variable (*cf.* section 4.3.3), only corrections for the numerator are calculated since the impact on the denominator is negligible. Since ρ represents a transverse energy density, the variable V_{raw} studied here is therefore $H/E \times E_T$. Its distribution is drawn as a function of E_T in figure C.3. The corresponding effective area is shown in figure C.4. A strong dependency is seen as a function of the supercluster pseudorapidity. A fit to the parametric function:

$$E_A(\eta) = A \times \cosh(B\eta) + C$$

is therefore performed. The results of the fits are given in the figure C.4.

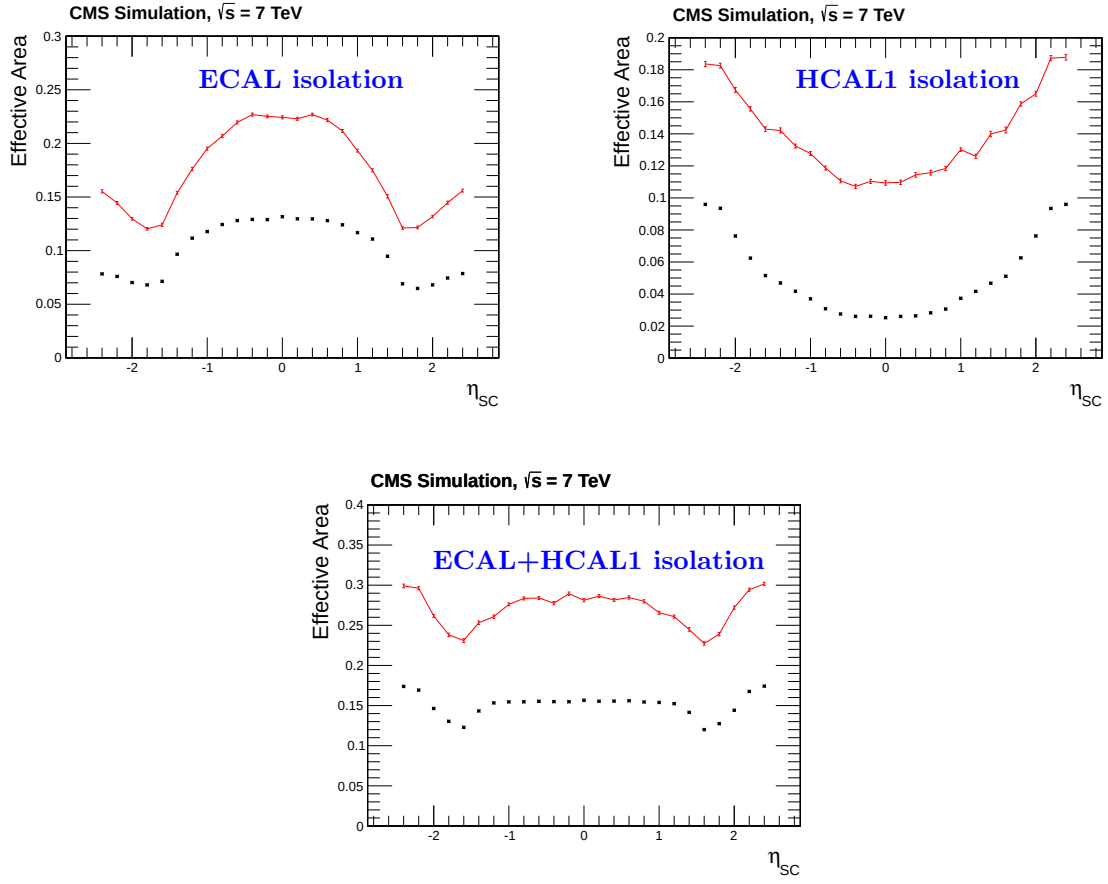


Figure C.2: The effective areas for the ECAL and HCAL1 isolation variables separately (top) and together (bottom) as a function of the supercluster pseudorapidity obtained using the slopes of the mean isolation value (black) or the 95% efficient cut curve (red).

C. EFFECTIVE AREAS MEASUREMENT AT $\sqrt{s} = 7$ TEV

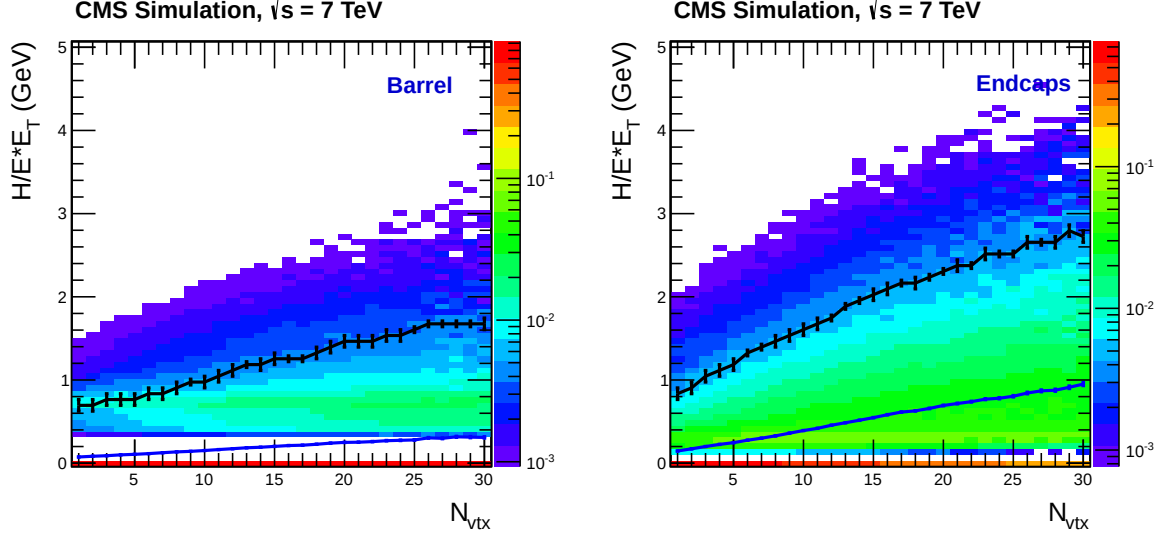


Figure C.3: The distribution of the $H_T = H/E \times E_T$ versus the number of primary vertices in the event for electrons with $E_T > 35$ GeV coming from a Drell-Yan simulated sample at $\sqrt{s} = 7$ TeV. Each slice in constant number of vertices is normalized to 1. For a given number of vertexes, the blue dots indicate the mean value of the projected distribution while the black dots indicate the value below which 95% of the events are contained.

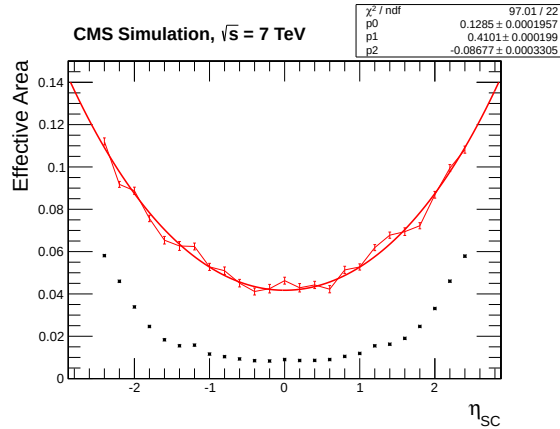


Figure C.4: The Effective Area for the $H_T = H/E \times E_T$ variable as a function of the supercluster pseudorapidity obtained using the slopes of the mean $H_T = H/E \times E_T$ value (black) or the 95% efficient cut curve (red). The results of a fit to the parametric function $p_0 \times \cosh(p_1 \eta) + p_2$ are also shown.

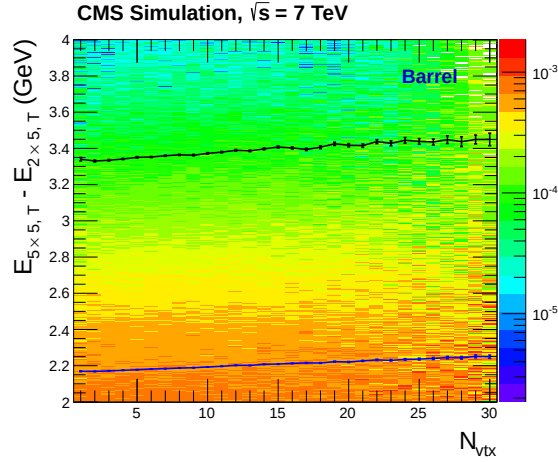


Figure C.5: The distribution of the $E_{5 \times 5} - E_{2 \times 5}$ versus the number of primary vertices in the event for electrons with $E_T > 35$ GeV coming from a Drell-Yan simulated sample at $\sqrt{s} = 7$ TeV. Each slice in constant number of vertices is normalized to 1. For a given number of vertexes, the blue dots indicate the mean value of the projected distribution while the black dots indicate the value below which 95% of the events are contained.

C.4 Transverse shower shape cut correction

For GSF electron candidates reconstructed in the barrel, a cut on the transverse shower shape is applied. It is based on the $E_{1 \times 5}$, $E_{2 \times 5}$ and $E_{5 \times 5}$ variables (*cf.* section 4.3.4). Because those variables include most of the electron energy which is much larger than the pile-up contribution, applying the procedure described above to measure the corresponding effective areas leads to large uncertainties. The effective area is therefore rather measured for the variable $E_{5 \times 5} - E_{2 \times 5}$ and then rescaled by a geometric factor of 1/3, 2/3 and 5/3 for $E_{1 \times 5}$, $E_{2 \times 5}$ and $E_{5 \times 5}$ respectively. The evolution of $E_{5 \times 5} - E_{2 \times 5}$ with the number of reconstructed vertices in the event is shown in figure C.5. The corresponding effective area is shown in figure C.6. Since the effective area shows no dependency with η , a straight line is fitted. An average value of 0.0055 is found. This corresponds to effective areas of 0.0018, 0.0037 and 0.0092 for $E_{1 \times 5,T}$, $E_{2 \times 5,T}$ and $E_{5 \times 5,T}$ respectively.

C. EFFECTIVE AREAS MEASUREMENT AT $\sqrt{s} = 7$ TEV

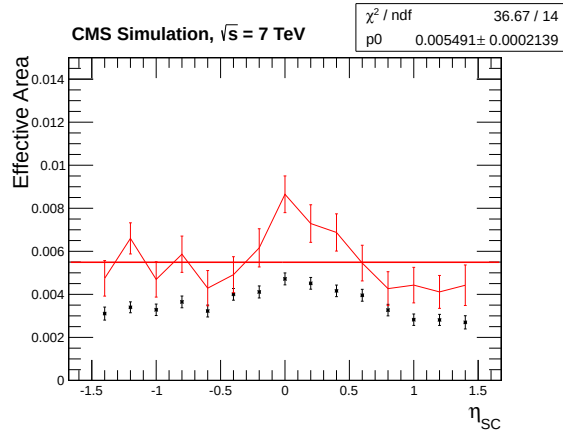


Figure C.6: The Effective Area for the $E_{5 \times 5} - E_{2 \times 5}$ variable as a function of the super-cluster pseudorapidity obtained using the slopes of the mean $E_{5 \times 5} - E_{2 \times 5}$ value (black) or the 95% efficient cut curve (red). The results of a fit to a straight line is also shown.

Appendix D

The Collins-Soper frame

This appendix derives the expression of the $\cos \theta_{CS}$ variable given in equation 7.4 of chapter 7.

D.1 Definition

When the frontal collision of two protons leads to the fusion of a $q\bar{q}$ pair into a photon or a Z boson decaying into a $l\bar{l}$ pair, an interesting quantity to study is the angle θ between the direction of the negative lepton and the quark in the rest frame of the $l\bar{l}$ pair. This variable is however not directly accessible in a hadron collider since the $q\bar{q}$ pair never occurs exactly in the direction of the pp collision: the radiation of gluons by the quark and antiquark before they interact will indeed induce a non zero transverse momentum of the pair.

In a paper [134], Collins and Soper suggested to use a variable θ_{CS} , which can be calculated using only quantities related to the dilepton pair and which approximates θ in case of a low transverse momentum for the pair.

Let $\vec{p}_A$ and $\vec{p}_B$ be the 3-momenta of the two colliding hadrons in the laboratory frame where A is the hadron providing the quark and B the hadron providing the antiquark, and let $\vec{p}'_A$ and $\vec{p}'_B$ be their values in a frame in the center of mass of the dilepton pair. The frame used to calculate θ_{CS} (CS frame) is defined as the frame in the center of mass of the dilepton pair for which:

- the z -axis bisects $\vec{p}'_A$ and $-\vec{p}'_B$.
- the x -axis is orthogonal to the z -axis, lies also in the plane spanned by $\vec{p}'_A$ and $\vec{p}'_B$ and its direction is chosen such that $p'_{x,A}$ and $p'_{x,B}$ are negative.

One then defines θ_{CS} as the angle between the negative lepton momentum and the z -axis.

D.2 Calculation of $\cos \theta_{CS}$

The following calculation apply for head-on collisions of equal energy protons. Let $\mathcal{R}$ be the laboratory frame in which the direction of the proton coincides with the z -axis. One

D. THE COLLINS-SOPER FRAME

defines the x -axis such as it coincides with the direction of the $q\bar{q}$ pair in the transverse plane and the y -axis as the direction perpendicular to the x and z -axis. Let $\mathcal{R}'$ be the Collins-Soper frame.

The coordinates of a quadrivector (E', p'_x, p'_y, p'_z) in the Collins-Soper frame is obtained from its coordinates in the laboratory frame (E, p_x, p_y, p_z) by applying a boost of parameters $(\beta_x, 0, \beta_z)$ followed by a rotation of parameter α in the (x, z) plane. Namely:

$$\begin{pmatrix} E' \\ p'_x \\ p'_y \\ p'_z \end{pmatrix} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & \cos \alpha & 0 & \sin \alpha \\ 0 & 0 & 1 & 0 \\ 0 & -\sin \alpha & 0 & \cos \alpha \end{pmatrix} \begin{pmatrix} \gamma & -\beta_x \gamma & 0 & -\beta_z \gamma \\ -\beta_x \gamma & 1 + (\gamma - 1) \frac{\beta_x^2}{\beta^2} & 0 & (\gamma - 1) \frac{\beta_x \beta_z}{\beta^2} \\ 0 & 0 & 1 & 0 \\ -\beta_z \gamma & (\gamma - 1) \frac{\beta_x \beta_z}{\beta^2} & 0 & 1 + (\gamma - 1) \frac{\beta_z^2}{\beta^2} \end{pmatrix} \begin{pmatrix} E \\ p_x \\ p_y \\ p_z \end{pmatrix} \quad (\text{D.1})$$

where $\beta = \sqrt{\beta_x^2 + \beta_z^2}$ and $\gamma = \frac{1}{\sqrt{1 - \beta^2}}$. The product of the boost and rotation matrices is:

$$\begin{pmatrix} \gamma & -\beta_x \gamma & 0 & -\beta_z \gamma \\ -\cos \alpha \beta_x \gamma - \sin \alpha \beta_z \gamma & \cos \alpha \left(1 + \beta_x^2 \frac{\gamma^2}{\gamma^2 + 1} \right) + \sin \alpha \beta_x \beta_z \frac{\gamma^2}{\gamma^2 + 1} & 0 & \cos \alpha \beta_x \beta_z \frac{\gamma^2}{\gamma^2 + 1} + \sin \alpha \left(1 + \beta_z^2 \frac{\gamma^2}{\gamma^2 + 1} \right) \\ 0 & 0 & 1 & 0 \\ \sin \alpha \beta_x \gamma - \cos \alpha \beta_z \gamma & -\sin \alpha \left(1 + \beta_x^2 \frac{\gamma^2}{\gamma^2 + 1} \right) + \cos \alpha \beta_x \beta_z \frac{\gamma^2}{\gamma^2 + 1} & 0 & -\sin \alpha \beta_x \beta_z \frac{\gamma^2}{\gamma^2 + 1} + \cos \alpha \left(1 + \beta_z^2 \frac{\gamma^2}{\gamma^2 + 1} \right) \end{pmatrix} \quad (\text{D.2})$$

using the fact that $\frac{\gamma - 1}{\beta^2} = \frac{\gamma^2}{\gamma + 1}$.

The constraint on α comes from the fact that the direction of two beams must cross the z' -axis with the same angle. For head on collision of equal energy protons, one has:

$$\begin{cases} p_A = (E_A, 0, 0, p_{z,A}) \\ p_B = (E_B, 0, 0, p_{z,B}) \\ p'_A = (E'_A, p'_{x,A}, 0, p'_{z,A}) \\ p'_B = (E'_B, p'_{x,B}, 0, p'_{z,B}) \end{cases} \quad (\text{D.3})$$

and the condition is:

$$\frac{E'_1}{p'_{x,1}} = \frac{E'_2}{p'_{x,2}} \quad (\text{D.4})$$

or in terms of coordinates in the laboratory frame:

$$\begin{aligned} & \frac{\gamma E_A - \beta_z \gamma E_A}{(-\cos \alpha \beta_x - \sin \alpha \beta_z \gamma) E_A + \left(\cos \alpha \beta_x \beta_z \frac{\gamma^2}{\gamma^2 + 1} + \sin \alpha (1 + \beta_z^2 \frac{\gamma^2}{\gamma^2 + 1}) \right) E_A} \\ &= \frac{\gamma E_B + \beta_z \gamma E_B}{(-\cos \alpha \beta_x - \sin \alpha \beta_z \gamma) E_B - \left(\cos \alpha \beta_x \beta_z \frac{\gamma^2}{\gamma^2 + 1} + \sin \alpha (1 + \beta_z^2 \frac{\gamma^2}{\gamma^2 + 1}) \right) E_B} \end{aligned} \quad (\text{D.5})$$

which leads to the conditions:

$$\begin{cases} \sin \alpha = \frac{\beta_x \beta_z \frac{\gamma}{\gamma + 1}}{\sqrt{1 - \beta_z^2}} \\ \cos \alpha = \frac{1 - \beta_z^2 \frac{\gamma}{\gamma + 1}}{\sqrt{1 - \beta_z^2}} \end{cases} \quad (\text{D.6})$$

Denoting $p_{l\bar{l}} = (E_{l\bar{l}}, p_{x,l\bar{l}}, 0, p_{z,l\bar{l}})$ the momentum of the dilepton pair in the laboratory frame and $M_{l\bar{l}}$ its mass, the boost parameters are given by:

$$\begin{cases} \beta_x = \frac{p_{x,l\bar{l}}}{E_{l\bar{l}}} \\ \beta_z = \frac{p_{z,l\bar{l}}}{E_{l\bar{l}}} \\ \gamma = \frac{E_{l\bar{l}}}{M_{l\bar{l}}} \end{cases} \quad (\text{D.7})$$

and $\cos \alpha$ and $\sin \alpha$ can be reexpressed as:

$$\begin{cases} \sin \alpha = \frac{p_{x,l\bar{l}} p_{z,l\bar{l}}}{(E_{l\bar{l}} + M_{l\bar{l}}) \sqrt{p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2}} \\ \cos \alpha = \frac{p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2 + M_{l\bar{l}} E_{l\bar{l}}}{(E_{l\bar{l}} + M_{l\bar{l}}) \sqrt{p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2}} \end{cases} \quad (\text{D.8})$$

The variable $\cos \theta_{CS}$ is by definition the ratio of the longitudinal momentum of the lepton $p'_{z,l}$ over its energy E'_l in the CS frame:

$$\cos \theta_{CS} = \frac{p'_{z,l}}{E'_l} \quad (\text{D.9})$$

In the CS frame the electron and positron energies are equal and therefore $E'_l = M_{l\bar{l}}/2$ and $p'_{z,l}$ has the following expression in terms of variables in the laboratory frame:

$$\begin{aligned} p'_{z,l} = & E_l \cdot \left(\frac{p_{x,l\bar{l}}^2 p_{z,l\bar{l}}}{M_{l\bar{l}}(E_{l\bar{l}} + M_{l\bar{l}}) \sqrt{p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2}} - \frac{p_{z,l\bar{l}}(p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2 + M_{l\bar{l}} E_{l\bar{l}})}{M_{l\bar{l}}(E_{l\bar{l}} + M_{l\bar{l}}) \sqrt{p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2}} \right) \\ & + p_{x,l} \cdot \left(-\frac{p_{x,l\bar{l}} p_{z,l\bar{l}}(p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2 + M_{l\bar{l}} E_{l\bar{l}})}{M_{l\bar{l}}(E_{l\bar{l}} + M_{l\bar{l}})^2 \sqrt{p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2}} + \frac{p_{x,l\bar{l}} p_{z,l\bar{l}}(p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2 + M_{l\bar{l}} E_{l\bar{l}})}{M_{l\bar{l}}(E_{l\bar{l}} + M_{l\bar{l}})^2 \sqrt{p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2}} \right) \\ & + p_{y,l} \cdot 0 \\ & + p_{z,l} \cdot \left(\frac{-p_{x,l\bar{l}}^2 p_{z,l\bar{l}}}{M_{l\bar{l}}(E_{l\bar{l}} + M_{l\bar{l}})^2 \sqrt{p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2}} + \frac{(p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2 + M_{l\bar{l}} E_{l\bar{l}})(p_{z,l\bar{l}}^2 + M_{l\bar{l}}^2 + M_{l\bar{l}} E_{l\bar{l}})}{M_{l\bar{l}}(E_{l\bar{l}} + M_{l\bar{l}})^2 \sqrt{p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2}} \right) \end{aligned} \quad (\text{D.10})$$

One has $E_{l\bar{l}} = E_l + E_{\bar{l}}$ and $p_{z,l\bar{l}} = p_{z,l} + p_{z,\bar{l}}$, where $p_{z,l}$, E_l ($p_{z,\bar{l}}$, $E_{\bar{l}}$) are the longitudinal impulsions and energy in the laboratory frame of the lepton and antilepton respectively, and equation D.10 thus reduces to:

$$p'_{z,l} = \frac{p_{z,l} E_{\bar{l}} - p_{z,\bar{l}} E_l}{\sqrt{p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2}} \quad (\text{D.11})$$

which leads to the announced result for $\cos \theta_{CS}$:

$$\cos \theta_{CS} = \frac{2}{M_{l\bar{l}}} \frac{(p_{z,l} E_{\bar{l}} - p_{z,\bar{l}} E_l)}{\sqrt{p_{x,l\bar{l}}^2 + M_{l\bar{l}}^2}} \quad (\text{D.12})$$

D. THE COLLINS-SOPER FRAME

References

- [1] The ATLAS Collaboration. “Observation of a new particle in the search for the Standard Model Higgs boson with the ATLAS detector at the LHC”. *Phys.Lett.*, B716:1–29, 2012. doi: 10.1016/j.physletb.2012.08.020. [1](#)
- [2] The CMS Collaboration. “Observation of a new boson at a mass of 125 GeV with the CMS experiment at the LHC”. *Phys.Lett.*, B716:30–61, 2012. doi: 10.1016/j.physletb.2012.08.021. [1](#)
- [3] The CMS Collaboration. “Search for Resonances in the Dilepton Mass Distribution in pp Collisions at $\sqrt{s} = 8$ TeV”. *CMS-PAS-EXO-12-061*, 2013. [2](#), [145](#)
- [4] The CMS Collaboration. “Searches for evidence of physics beyond the standard model in dilepton mass spectra in pp collisions at $\sqrt{s} = 8$ TeV”. *In preparation*. [2](#), [145](#), [173](#)
- [5] The CMS Collaboration. “Search for Resonances in the Dilepton Mass Distribution in *pp* Collisions at $\sqrt{s} = 7$ TeV”. *JHEP*, 1105:093, 2011. doi: 10.1007/JHEP05(2011)093. [2](#), [145](#)
- [6] The CMS Collaboration. “Search for narrow resonances in dilepton mass spectra in *pp* collisions at $\sqrt{s} = 7$ TeV”. *Phys.Lett.*, B714:158–179, 2012. doi: 10.1016/j.physletb.2012.06.051. [2](#)
- [7] The CMS Collaboration. “Search for heavy narrow dilepton resonances in *pp* collisions at $\sqrt{s} = 7$ TeV and $\sqrt{s} = 8$ TeV”. *Phys.Lett.*, B720:63–82, 2013. doi: 10.1016/j.physletb.2013.02.003. [2](#), [145](#), [173](#)
- [8] The CMS Collaboration. “CMS 2013 Public Electron Performance Results”. <https://twiki.cern.ch/twiki/bin/view/CMSPublic/EGMElectronsMoriond2013>, 2013. [2](#), [45](#)
- [9] The CMS Collaboration. “Performance of electron reconstruction and selection with the CMS detector at $\sqrt{s} = 8$ TeV”. *In preparation*, 2014. [2](#), [72](#), [140](#)
- [10] Michael E. Peskin and Daniel V. Schroeder. “*An Introduction to quantum field theory*”. 1995. ISBN 9780201503975. [3](#)
- [11] S.L. Glashow. “Partial Symmetries of Weak Interactions”. *Nucl.Phys.*, 22:579–588, 1961. doi: 10.1016/0029-5582(61)90469-2. [3](#)

REFERENCES

- [12] Steven Weinberg. “A Model of Leptons”. *Phys.Rev.Lett.*, 19:1264–1266, 1967. doi: 10.1103/PhysRevLett.19.1264. [3](#)
- [13] Abdus Salam and John Clive Ward. “Electromagnetic and weak interactions”. *Phys.Lett.*, 13:168–171, 1964. doi: 10.1016/0031-9163(64)90711-5. [3](#)
- [14] J. Beringer et al. “Review of Particle Physics (RPP)”. *Phys.Rev.*, D86:010001, 2012. doi: 10.1103/PhysRevD.86.010001. [4](#), [9](#), [15](#), [18](#), [165](#)
- [15] Noether Emmy. “invariante variationsprobleme”. *Nachr. v. d. Ges. d. Wiss. zu Gttingen*, pages 235–259, 1918. [4](#)
- [16] F. Englert and R. Brout. “Broken Symmetry and the Mass of Gauge Vector Mesons”. *Phys.Rev.Lett.*, 13:321–323, 1964. doi: 10.1103/PhysRevLett.13.321. [8](#)
- [17] Peter W. Higgs. “Broken symmetries, massless particles and gauge fields”. *Phys.Lett.*, 12:132–133, 1964. doi: 10.1016/0031-9163(64)91136-9. [8](#)
- [18] Peter W. Higgs. “Broken Symmetries and the Masses of Gauge Bosons”. *Phys.Rev.Lett.*, 13:508–509, 1964. doi: 10.1103/PhysRevLett.13.508. [8](#)
- [19] Gerard ’t Hooft and M.J.G. Veltman. “Regularization and Renormalization of Gauge Fields”. *Nucl.Phys.*, B44:189–213, 1972. doi: 10.1016/0550-3213(72)90279-9. [9](#)
- [20] M. Peskin. “Beyond the Standard Model”. *Lectures notes of the 1996 European School of High-Energy Physics, Carry-le-Rouet, France*, 1996. [10](#)
- [21] David J. Gross and Frank Wilczek. “Ultraviolet Behavior of Nonabelian Gauge Theories”. *Phys.Rev.Lett.*, 30:1343–1346, 1973. doi: 10.1103/PhysRevLett.30.1343. [11](#)
- [22] H. David Politzer. “Reliable Perturbative Results for Strong Interactions?”. *Phys.Rev.Lett.*, 30:1346–1349, 1973. doi: 10.1103/PhysRevLett.30.1346. [11](#)
- [23] S.D. Drell and Tung-Mow Yan. “Massive Lepton Pair Production in Hadron-Hadron Collisions at High-Energies”. *Phys.Rev.Lett.*, 25:316–320, 1970. doi: 10.1103/PhysRevLett.25.316. [11](#)
- [24] Jonathan L. Rosner. “Forward - backward asymmetries in hadronically produced lepton pairs”. *Phys.Rev.*, D54:1078–1082, 1996. doi: 10.1103/PhysRevD.54.1078. [12](#), [19](#), [20](#)
- [25] Stephen P. Martin. “A Supersymmetry primer”. *Adv.Ser.Direct.High Energy Phys.*, 21:1–153, 2010. doi: 10.1142/9789814307505_0001. [14](#), [16](#)
- [26] Super-Kamiokande Collaboration. “Evidence for oscillation of atmospheric neutrinos”. *Phys.Rev.Lett.*, 81:1562–1567, 1998. doi: 10.1103/PhysRevLett.81.1562. [15](#)

-
- [27] DOUBLE-CHOOZ Collaboration. “Indication for the disappearance of reactor electron antineutrinos in the Double Chooz experiment”. *Phys.Rev.Lett.*, 108:131801, 2012. doi: 10.1103/PhysRevLett.108.131801. [15](#)
- [28] Vera C. Rubin and Jr. Ford, W. Kent. “Rotation of the Andromeda Nebula from a Spectroscopic Survey of Emission Regions”. *Astrophys.J.*, 159:379–403, 1970. doi: 10.1086/150317. [15](#)
- [29] Katherine Freese. “Review of Observational Evidence for Dark Matter in the Universe and in upcoming searches for Dark Stars”. *astro-ph/0812.4005*, 2008. [15](#)
- [30] Adam G. Riess et al. “Observational evidence from supernovae for an accelerating universe and a cosmological constant”. *Astron.J.*, 116:1009–1038, 1998. doi: 10.1086/300499. [15](#)
- [31] S. Perlmutter et al. “Measurements of Omega and Lambda from 42 high redshift supernovae”. *Astrophys.J.*, 517:565–586, 1999. doi: 10.1086/307221. [15](#)
- [32] Paul Langacker. “The Physics of Heavy Z' Gauge Bosons”. *Rev.Mod.Phys.*, 81:1199–1228, 2009. doi: 10.1103/RevModPhys.81.1199. [15](#)
- [33] R parity Working Group Collaboration. “Searching for R parity violation at Run II of the Tevatron”. *hep-ph/9906224*, 1999. [16](#)
- [34] H. Georgi and S.L. Glashow. “Unity of All Elementary Particle Forces”. *Phys.Rev.Lett.*, 32:438–441, 1974. doi: 10.1103/PhysRevLett.32.438. [17](#)
- [35] Michael E. Peskin. “Beyond the standard model”. *hep-ph/9705479*, 1997. [18](#)
- [36] Edward Witten. “Symmetry Breaking Patterns in Superstring Models”. *Nucl.Phys.*, B258:75, 1985. doi: 10.1016/0550-3213(85)90603-0. [19](#)
- [37] Nima Arkani-Hamed, Savas Dimopoulos, and G.R. Dvali. “The Hierarchy problem and new dimensions at a millimeter”. *Phys.Lett.*, B429:263–272, 1998. doi: 10.1016/S0370-2693(98)00466-3. [20](#)
- [38] Lisa Randall and Raman Sundrum. “A Large mass hierarchy from a small extra dimension”. *Phys.Rev.Lett.*, 83:3370–3373, 1999. doi: 10.1103/PhysRevLett.83.3370. [21](#)
- [39] The DELPHI Collaboration. “A Study of radiative muon pair events at Z^0 energies and limits on an additional Z' gauge boson”. *Z.Phys.*, C65:603–618, 1995. doi: 10.1007/BF01578669. [23](#)
- [40] Ross Diener, Stephen Godfrey, and Ismail Turan. “Constraining Extra Neutral Gauge Bosons with Atomic Parity Violation Measurements”. *Phys.Rev.*, D86:115017, 2012. doi: 10.1103/PhysRevD.86.115017. [23](#)

REFERENCES

- [41] Jens Erler, Paul Langacker, Shoaib Munir, and Eduardo Rojas. “Improved Constraints on Z-prime Bosons from Electroweak Precision Data”. *JHEP*, 0908:017, 2009. doi: 10.1088/1126-6708/2009/08/017. [23](#)
- [42] L3 Collaboration OPAL Collaboration LEP Electroweak Working Group ALEPH Collaboration, DELPHI Collaboration. “A Combination of preliminary electroweak measurements and constraints on the standard model”. 2006. [23](#)
- [43] The CDF Collaboration. “Search for High-Mass e^+e^- Resonances in $p\bar{p}$ Collisions at $\sqrt{s} = 1.96$ TeV”. *Phys.Rev.Lett.*, 102:031801, 2009. doi: 10.1103/PhysRevLett.102.031801. [23](#), [145](#)
- [44] The D0 Collaboration. “Search for a heavy neutral gauge boson in the dielectron channel with 5.4 fb⁻¹ of $p\bar{p}$ collisions at $\sqrt{s} = 1.96$ TeV”. *Phys.Lett.*, B695:88–94, 2011. doi: 10.1016/j.physletb.2010.10.059. [23](#), [145](#)
- [45] The CDF Collaboration. “Search for New Dielectron Resonances and Randall-Sundrum Gravitons at the Collider Detector at Fermilab”. *Phys.Rev.Lett.*, 107:051801, 2011. doi: 10.1103/PhysRevLett.107.051801. [23](#), [145](#)
- [46] UA1 Collaboration. “Experimental Observation of Lepton Pairs of Invariant Mass Around 95-GeV/c² at the CERN SPS Collider”. *Phys.Lett.*, B126:398–410, 1983. doi: 10.1016/0370-2693(83)90188-0. [25](#)
- [47] UA2 Collaboration. “Evidence for $Z^0 \rightarrow e^+e^-$ at the CERN anti-p p Collider”. *Phys.Lett.*, B129:130–140, 1983. doi: 10.1016/0370-2693(83)90744-X. [25](#)
- [48] UA1 Collaboration. “Experimental Observation of Events with Large Missing Transverse Energy Accompanied by a Jet Or a Photon(s) in p anti-p Collisions at $\sqrt{s}=540$ GeV”. *Phys.Lett.*, B139:115, 1984. doi: 10.1016/0370-2693(84)90046-7. [25](#)
- [49] UA2 Collaboration. “Observation of Single Isolated Electrons of High Transverse Momentum in Events with Missing Transverse Energy at the CERN anti-p p Collider”. *Phys.Lett.*, B122:476–485, 1983. doi: 10.1016/0370-2693(83)91605-2. [25](#)
- [50] LEP Electroweak. “Electroweak Measurements in Electron-Positron Collisions at W-Boson-Pair Energies at LEP”. *Phys.Rept.*, 532:119–244, 2013. doi: 10.1016/j.physrep.2013.07.004. [25](#)
- [51] ALEPH Collaboration et al. “Precision electroweak measurements on the Z resonance”. *Phys.Rept.*, 427:257–454, 2006. doi: 10.1016/j.physrep.2005.12.006. [25](#)
- [52] The D0 Collaboration. “Search for high mass top quark production in $p\bar{p}$ collisions at $\sqrt{s} = 1.8$ TeV”. *Phys.Rev.Lett.*, 74:2422–2426, 1995. doi: 10.1103/PhysRevLett.74.2422. [25](#)
- [53] The CDF Collaboration. “Observation of top quark production in $p\bar{p}$ collisions”. *Phys.Rev.Lett.*, 74:2626–2631, 1995. doi: 10.1103/PhysRevLett.74.2626. [25](#)

-
- [54] Lyndon Evans and Philip Bryant. “LHC Machine”. *JINST*, 3:S08001, 2008. doi: 10.1088/1748-0221/3/08/S08001. 25
- [55] The ATLAS Collaboration. “The ATLAS Experiment at the CERN Large Hadron Collider”. *JINST*, 3:S08003, 2008. doi: 10.1088/1748-0221/3/08/S08003. 25
- [56] The ALICE Collaboration. “The ALICE experiment at the CERN LHC”. *JINST*, 3:S08002, 2008. doi: 10.1088/1748-0221/3/08/S08002. 25
- [57] The CMS Collaboration. “The CMS experiment at the CERN LHC”. *JINST*, 3:S08004, 2008. doi: 10.1088/1748-0221/3/08/S08004. 25, 49
- [58] The LHCb Collaboration. “The LHCb Detector at the LHC”. *JINST*, 3:S08005, 2008. doi: 10.1088/1748-0221/3/08/S08005. 25
- [59] W. Herr and B. Muratori. “Concept of luminosity”. *CAS - CERN Accelerator School: Intermediate Course on Accelerator Physics*, 2003. 28
- [60] The CMS Collaboration. “Measurement of CMS Luminosity”. *CMS-PAS-EWK-10-004*, 2010. 28
- [61] The CMS Collaboration. “CMS Luminosity Based on Pixel Cluster Counting - Summer 2013 Update”. *CMS-PAS-LUM-13-001*, 2013. 29, 146
- [62] Cern announces lhc restart schedule. “<http://home.web.cern.ch/about/updates/2014/06/cern-announces-lhc-restart-schedule>”. 29
- [63] P. Wells on behalf of ATLAS and CMS. “Physics Program of the LHC Upgrades (talk at the Lepton Photon 2013 Conference)”. <http://indico.cern.ch/event/242095/session/7/contribution/14/material/slides/0.pdf>. 31
- [64] The CMS Collaboration. “CMS Tracker Detector Performance Results”. <https://twiki.cern.ch/twiki/bin/view/CMSPublic/DPGResultsTRK>, 2012. 35
- [65] The CMS Collaboration. “Alignment of the CMS tracker with LHC and cosmic ray data”. *JINST*, 9:P06009, 2014. doi: 10.1088/1748-0221/9/06/P06009. 35
- [66] P. Adzic, R. Alemany-Fernandez, C.B. Almeida, N.M. Almeida, G. Anagnostou, et al. “Energy resolution of the barrel of the CMS electromagnetic calorimeter”. *JINST*, 2:P04004, 2007. doi: 10.1088/1748-0221/2/04/P04004. 36
- [67] The CMS Collaboration. “Energy Calibration and Resolution of the CMS Electromagnetic Calorimeter in pp Collisions at $\sqrt{s} = 7$ TeV”. *JINST*, 8:P09009, 2013. doi: 10.1088/1748-0221/8/09/P09009. 37, 43
- [68] The CMS Collaboration. “ECAL DPG Approved Results”. <https://twiki.cern.ch/twiki/bin/view/CMSPublic/EcalDPGResults>, 2012. 37, 38, 40, 41, 44

REFERENCES

- [69] W. Adam, B. Mangano, T. Speer, and T. Todorov. “Track reconstruction in the CMS tracker”. *CMS-NOTE-2006-041*, 2005. [49](#)
- [70] S. Cucciarelli, M. Konecki, D. Kotlinski, and T. Todorov. “Track reconstruction, primary vertex finding and seed generation with the pixel detector”. *CMS-NOTE-2006-026*, 2006. [49](#)
- [71] The CMS Collaboration. “Particle-Flow Event Reconstruction in CMS and Performance for Jets, Taus, and MET”. *CMS-PAS-PFT-09-001*, 2009. [50](#)
- [72] Florian Beaudette. “The CMS Particle Flow Algorithm”. *hep-ex/1401.8155*, pages 295–304, 2014. [50](#)
- [73] Matteo Cacciari, Gavin P. Salam, and Gregory Soyez. “The Anti-k(t) jet clustering algorithm”. *JHEP*, 0804:063, 2008. doi: 10.1088/1126-6708/2008/04/063. [50](#)
- [74] The CMS Collaboration. “CMS Particle Flow and Tau Identification Results”. <https://twiki.cern.ch/twiki/bin/view/CMSPublic/PhysicsResultsPFT>, 2014. [54](#)
- [75] The CMS Collaboration. “Search for high mass resonances decaying into τ lepton pairs in pp collisions at $\sqrt{s} = 7$ TeV”. *Phys.Lett.*, B716:82–102, 2012. doi: 10.1016/j.physletb.2012.07.062. [53](#)
- [76] The CMS Collaboration. “Performance of Missing Transverse Momentum Reconstruction Algorithms in Proton-Proton Collisions at $\sqrt{s} = 8$ TeV with the CMS Detector”. *CMS-PAS-JME-12-002*, 2012. [56](#)
- [77] James Stirling. “Parton luminosity and cross section plots”. <http://www.hep.ph.ic.ac.uk/~wstirlin/plots/plots.html>, 2014. [58](#)
- [78] The TOTEM Collaboration. “Luminosity-Independent Measurement of the Proton-Proton Total Cross Section at $\sqrt{s} = 8$ TeV”. *Phys.Rev.Lett.*, 111(1):012001, 2013. doi: 10.1103/PhysRevLett.111.012001. [57](#)
- [79] The H1 Collaboration. “Inclusive Deep Inelastic Scattering at High Q^2 with Longitudinally Polarised Lepton Beams at HERA”. *JHEP*, 1209:061, 2012. doi: 10.1007/JHEP09(2012)061. [60](#)
- [80] The ZEUS Collaboration. “Measurement of high- Q^2 neutral current deep inelastic e^+p scattering cross sections with a longitudinally polarized positron beam at HERA”. *Phys.Rev.*, D87(5):052014, 2013. doi: 10.1103/PhysRevD.87.052014. [60](#)
- [81] The CMS Collaboration. “Measurements of the differential and double-differential Drell-Yan cross sections in proton-proton collisions at 8 TeV”. *CMS-PAS-SMP-14-003*, 2014. [60](#)
- [82] The CMS Collaboration. “Measurement of the muon charge asymmetry in inclusive $pp \rightarrow W + X$ production at $\sqrt{s}=7$ TeV and an improved determination of light parton distribution functions”. *CMS-SMP-12-021*, 2013. [60](#)

-
- [83] Yuri L. Dokshitzer. “Calculation of the Structure Functions for Deep Inelastic Scattering and e^+e^- Annihilation by Perturbation Theory in Quantum Chromodynamics”. *Sov.Phys.JETP*, 46:641–653, 1977. [60](#)
- [84] Guido Altarelli and G. Parisi. “Asymptotic Freedom in Parton Language”. *Nucl.Phys.*, B126:298, 1977. doi: 10.1016/0550-3213(77)90384-4. [60](#)
- [85] V.N. Gribov and L.N. Lipatov. “Deep inelastic $e p$ scattering in perturbation theory”. *Sov.J.Nucl.Phys.*, 15:438–450, 1972. [60](#)
- [86] J. Pumplin, D.R. Stump, J. Huston, H.L. Lai, Pavel M. Nadolsky, et al. “New generation of parton distributions with uncertainties from global QCD analysis”. *JHEP*, 0207:012, 2002. doi: 10.1088/1126-6708/2002/07/012. [61](#)
- [87] Hung-Liang Lai, Marco Guzzi, Joey Huston, Zhao Li, Pavel M. Nadolsky, et al. “New parton distributions for collider physics”. *Phys.Rev.*, D82:074024, 2010. doi: 10.1103/PhysRevD.82.074024. [61](#)
- [88] A.D. Martin, W.J. Stirling, R.S. Thorne, and G. Watt. “Parton distributions for the LHC”. *Eur.Phys.J.*, C63:189–285, 2009. doi: 10.1140/epjc/s10052-009-1072-5. [61](#)
- [89] Richard D. Ball, Luigi Del Debbio, Stefano Forte, Alberto Guffanti, Jose I. Latorre, et al. “A first unbiased global NLO determination of parton distributions and their uncertainties”. *Nucl.Phys.*, B838:136–206, 2010. doi: 10.1016/j.nuclphysb.2010.05.008. [61](#)
- [90] The Durham HepData Project. “Online PDF plotting and calculation”. <http://hepdata.cedar.ac.uk/pdf/pdf3.htm>. [62](#)
- [91] John C. Collins, Davison E. Soper, and George F. Sterman. “Factorization of Hard Processes in QCD”. *Adv.Ser.Direct.High Energy Phys.*, 5:1–91, 1988. [63](#)
- [92] M.A. Dobbs, S. Frixione, Eric Laenen, K. Tollefson, H. Baer, et al. “Les Houches guidebook to Monte Carlo generators for hadron collider physics”. *hep-ph/0403045*, pages 411–459, 2004. [64](#)
- [93] Torbjorn Sjostrand, Stephen Mrenna, and Peter Z. Skands. “PYTHIA 6.4 Physics and Manual”. *JHEP*, 0605:026, 2006. doi: 10.1088/1126-6708/2006/05/026. [64](#)
- [94] Torbjorn Sjostrand, Stephen Mrenna, and Peter Z. Skands. “A Brief Introduction to PYTHIA 8.1”. *Comput.Phys.Commun.*, 178:852–867, 2008. doi: 10.1016/j.cpc.2008.01.036. [64](#), [177](#)
- [95] Johan Alwall, Michel Herquet, Fabio Maltoni, Olivier Mattelaer, and Tim Stelzer. “MadGraph 5 : Going Beyond”. *JHEP*, 1106:128, 2011. doi: 10.1007/JHEP06(2011)128. [64](#)

REFERENCES

- [96] Stefano Frixione, Paolo Nason, and Carlo Oleari. “Matching NLO QCD computations with Parton Shower simulations: the POWHEG method”. *JHEP*, 0711:070, 2007. doi: 10.1088/1126-6708/2007/11/070. [64](#), [154](#)
- [97] Ryan Gavin, Ye Li, Frank Petriello, and Seth Quackenbush. “FEWZ 2.0: A code for hadronic Z production at next-to-next-to-leading order”. *Comput.Phys.Commun.*, 182:2388–2403, 2011. doi: 10.1016/j.cpc.2011.06.008. [64](#), [67](#), [154](#), [155](#)
- [98] John C. Collins. “Sudakov form-factors”. *Adv.Ser.Direct.High Energy Phys.*, 5: 573–614, 1989. [65](#)
- [99] P. Bartalini, E.L. Berger, B. Blok, G. Calucci, R. Corke, et al. “Multi-Parton Interactions at the LHC”. *CMS-CR-2011-048*, 2011. [65](#)
- [100] The CMS Collaboration. “Measurement of the Underlying Event Activity at the LHC with $\sqrt{s} = 7$ TeV and Comparison with $\sqrt{s} = 0.9$ TeV”. *JHEP*, 1109:109, 2011. doi: 10.1007/JHEP09(2011)109. [65](#)
- [101] Bo Andersson, G. Gustafson, G. Ingelman, and T. Sjostrand. “Parton Fragmentation and String Dynamics”. *Phys.Rept.*, 97:31–145, 1983. doi: 10.1016/0370-1573(83)90080-7. [65](#)
- [102] John Allison, K. Amako, J. Apostolakis, H. Araujo, P.A. Dubois, et al. “Geant4 developments and applications. *IEEE Trans.Nucl.Sci.*, 53:270, 2006. doi: 10.1109/TNS.2006.869826. [66](#)
- [103] Micha Czakon, Paul Fiedler, and Alexander Mitov. “Total Top-Quark Pair-Production Cross Section at Hadron Colliders Through $O(\frac{4}{s})$ ”. *Phys.Rev.Lett.*, 110(25):252004, 2013. doi: 10.1103/PhysRevLett.110.252004. [67](#)
- [104] Nikolaos Kidonakis. “NNLL threshold resummation for top-pair and single-top production”. *Phys.Part.Nucl.*, 45(4):714–722, 2014. doi: 10.1134/S1063779614040091. [67](#)
- [105] J. M. Campbell, R. K. Ellis, and C. Williams. “MCFM - Monte Carlo for FeMtobarn processes”. <http://mcfm.fnal.gov/>, 2014. [67](#)
- [106] T. Binoth, J.P. Guillet, E. Pilon, and M. Werlen. “A Next-to-leading order study of photon pion and pion pair hadro production in the light of the Higgs boson search at the LHC”. *Eur.Phys.J.direct*, C4:7, 2002. doi: 10.1007/s1010502c0007. [67](#)
- [107] E. Meschi et al. “Electron Reconstruction in the CMS Electromagnetic Calorimeter”. *CMS-Note-2001/034*, 2001. [72](#)
- [108] W. Adam et al. “Reconstruction of Electrons with the Gaussian-Sum Filter in the CMS Tracker at the LHC”. *CMS-NOTE-2005-001*, 2005. [74](#)
- [109] R. Fruhwirth. “A Gaussian-mixture approximation of the BetheHeitler model of electron energy loss by bremsstrahlung”. *Computer Physics Communications*, 154, 2003. [74](#)

-
- [110] L. Thomas and B. Clerbaux. “High E_T electron selection presentation and optimisation in case of high pile-up conditions”. CMS AN-2012/258, 2012. [78](#), [98](#)
- [111] Matteo Cacciari and Gavin P. Salam. “Pileup subtraction using jet areas”. *Phys.Lett.*, B659:119–126, 2008. doi: 10.1016/j.physletb.2007.09.077. [98](#)
- [112] Matteo Cacciari, Gavin P. Salam, and Gregory Soyez. “The Catchment Area of Jets”. *JHEP*, 0804:005, 2008. doi: 10.1088/1126-6708/2008/04/005. [98](#)
- [113] Matteo Cacciari, Gavin P. Salam, and Gregory Soyez. “FastJet User Manual”. *Eur.Phys.J.*, C72:1896, 2012. doi: 10.1140/epjc/s10052-012-1896-2. [98](#)
- [114] L. Thomas and B. Clerbaux. “Measurement of the high energy electron selection efficiency using Drell-Yan events collected in 2012”. CMS AN-2013/174, 2013. [116](#)
- [115] The CMS Collaboration. “Measuring Electron Efficiencies at CMS with Early Data”. *CMS-PAS-EGM-07-001*, 2008. [116](#)
- [116] The CMS Collaboration. “Measurements of Inclusive W and Z Cross Sections in pp Collisions at $\sqrt{s} = 7$ TeV”. *JHEP*, 1101:080, 2011. doi: 10.1007/JHEP01(2011)080. [116](#)
- [117] The ATLAS Collaboration. “Measurement of the Z/γ^* boson transverse momentum distribution in pp collisions at $\sqrt{s} = 7$ TeV with the ATLAS detector”. *CERN-PH-EP-2014-075*, 2014. [120](#)
- [118] The CMS Collaboration. “Measurements of jet multiplicity and differential production cross sections of Z +jets events in proton-proton collisions at $\sqrt{s}=7$ TeV”. *CMS-SMP-12-017*, 2014. [120](#)
- [119] The CMS Collaboration. “Measurement of the double differential cross section of Z bosons produced in association with jets in pp collisions at $\sqrt{s} = 8$ TeV”. *CMS-PAS-SMP-14-009*, 2014. [120](#)
- [120] The ATLAS Collaboration. “Search for high-mass resonances decaying to dilepton final states in pp collisions at $\sqrt{s} = 7$ TeV with the ATLAS detector”. *JHEP*, 1211:138, 2012. doi: 10.1007/JHEP11(2012)138. [145](#)
- [121] The ATLAS Collaboration. “Search for high-mass dilepton resonances in pp collisions at $\sqrt{s} = 8$ TeV with the ATLAS detector”. *Phys.Rev.*, D (Accepted), 2014. [145](#), [174](#)
- [122] B. Clerbaux et al. “Update of the $Z' \rightarrow ee$ analysis with Jan 22 re-recoed data”. CMS AN-2013/359, 2013. [152](#)
- [123] S. Elgammal and B. Clerbaux. “Validation of ECAL calibration at high energy by the shower shape using Drell-Yan events in 2011 data”. CMS AN-2011/355, 2013. [152](#)

REFERENCES

- [124] Richard D. Ball et al. “Parton distributions with QED corrections”. *Nucl.Phys.*, B877:290–320, 2013. doi: 10.1016/j.nuclphysb.2013.10.010. [155](#)
- [125] G. Cowan. “*Statistical data analysis*”. 1998. ISBN 9780198501565. [165](#)
- [126] S. Schmitz et al. “Statistical inference in a Search for a Narrow Resonance”. CMS AN-2012/185, 2012. [165](#)
- [127] Michiel Botje, Jon Butterworth, Amanda Cooper-Sarkar, Albert de Roeck, Joel Feltess, et al. “The PDF4LHC Working Group Interim Recommendations”. *hep-ph/1101.0538*, 2011. [168](#)
- [128] M.R. Whalley, D. Bourilkov, and R.C. Group. “The Les Houches accord PDFs (LHAPDF) and LHAGLUE”. *hep-ph/0508110*, 2005. [168](#)
- [129] C.M. Carloni Calame, G. Montagna, O. Nicrosini, and A. Vicini. “Precision electroweak calculation of the production of a high transverse-momentum lepton pair at hadron colliders”. *JHEP*, 0710:109, 2007. doi: 10.1088/1126-6708/2007/10/109. [168](#)
- [130] Michal Czakon, Michelangelo L. Mangano, Alexander Mitov, and Juan Rojo. “Constraints on the gluon PDF from top quark pair production at hadron colliders”. *JHEP*, 1307:167, 2013. doi: 10.1007/JHEP07(2013)167. [168](#)
- [131] Elena Accomando, Diego Becciolini, Alexander Belyaev, Stefano Moretti, and Claire Shepherd-Themistocleous. “Z’ at the LHC: Interference and Finite Width Effects in Drell-Yan”. *JHEP*, 1310:153, 2013. doi: 10.1007/JHEP10(2013)153. [169](#)
- [132] R. Hamberg, W.L. van Neerven, and T. Matsuura. “A Complete calculation of the order $\alpha - s^2$ correction to the Drell-Yan K factor”. *Nucl.Phys.*, B359:343–405, 1991. doi: 10.1016/0550-3213(91)90064-5. [169](#)
- [133] The CMS Collaboration. “Projected Performance of an Upgraded CMS Detector at the LHC and HL-LHC: Contribution to the Snowmass Process”. 2013. [175](#), [176](#)
- [134] John C. Collins and Davison E. Soper. “Angular Distribution of Dileptons in High-Energy Hadron Collisions”. *Phys.Rev.*, D16:2219, 1977. doi: 10.1103/PhysRevD.16.2219. [180](#), [241](#)
- [135] L.D. Landau. “On the angular momentum of a two-photon system”. *Dokl.Akad.Nauk Ser.Fiz.*, 60:207–209, 1948. [184](#)
- [136] Chen-Ning Yang. “Selection Rules for the Dematerialization of a Particle Into Two Photons”. *Phys.Rev.*, 77:242–245, 1950. doi: 10.1103/PhysRev.77.242. [184](#)
- [137] Yanyan Gao, Andrei V. Gritsan, Zijin Guo, Kirill Melnikov, Markus Schulze, et al. “Spin determination of single-produced resonances at hadron colliders”. *Phys.Rev.*, D81:075022, 2010. doi: 10.1103/PhysRevD.81.075022. [184](#)

- [138] P. Osland, A.A. Pankov, N. Paver, and A.V. Tsytrinov. “Spin identification of the Randall-Sundrum resonance in lepton-pair production at the LHC”. *Phys.Rev.*, D78:035008, 2008. doi: 10.1103/PhysRevD.78.035008. [185](#)
- [139] G.L. Bayatian et al. “CMS technical design report, volume II: Physics performance”. *J.Phys.*, G34:995–1579, 2007. doi: 10.1088/0954-3899/34/6/S01. [185](#)
- [140] Alexander L. Read. “Presentation of search results: The CL(s) technique”. *J.Phys.*, G28:2693–2704, 2002. doi: 10.1088/0954-3899/28/10/313. [186](#)
- [141] The CMS Collaboration. “Study of the Mass and Spin-Parity of the Higgs Boson Candidate Via Its Decays to Z Boson Pairs”. *Phys.Rev.Lett.*, 110:081803, 2013. doi: 10.1103/PhysRevLett.110.081803. [186](#)