



Search for a SUSY signature through Vector Boson Fusion with
one muon, jets and missing transverse energy at $\sqrt{s} = 13$ TeV in
the CMS Experiment

Ph.D. Thesis

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Abstract

We present a search for supersymmetric (SUSY) electroweakino ($\tilde{\chi}_1^\pm, \tilde{\chi}_2^0$) production through Vector Boson Fusion (VBF) at $\sqrt{s} = 13$ TeV, carried out on proton-proton collision data corresponding to an integrated luminosity of 35.9 fb^{-1} collected by the CMS experiment in the 2016 LHC accelerator run. Our study is focused towards a final state composed by a single soft- p_T muon, large missing transverse momentum and two jets with large invariant mass and widely separated. The dijet system is the signature for the VBF topology, which is critical for background reduction. The event selection criteria were defined to maximize the SUSY signal significance under a compressed mass scenario where the wino-like electroweakino masses are nearly degenerate with respect to the bino-like Lightest SUSY particle ($\tilde{\chi}_1^0$). No data excess is observed over the estimated standard model backgrounds. Our results are combined with the ones from similar studies for the other lepton flavors (electron, tau and without any reconstructed lepton) to establish 95% CL exclusion limits for a slepton democratic scenario, assuming $m_{\tilde{\chi}_1^\pm} = m_{\tilde{\chi}_2^0}$ and slightly higher than $m_{\tilde{\chi}_1^0}$ by $\Delta m = m_{\tilde{\chi}_1^\pm} - m_{\tilde{\chi}_1^0} = 1$ GeV through 50 GeV range. Chargino masses are excluded up to $m_{\tilde{\chi}_1^\pm} < 300$ GeV (212 GeV) for $\Delta m = 30$ GeV (10 GeV).

Dedicated to the Creator of everything: Jesus.

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Introduction

The crucial piece of the puzzle to understand electroweak symmetry breaking through the Higgs mechanism, as predicted by the Standard Model (SM) [1–3], was finally discovered by both ATLAS and CMS experiments, at the LHC accelerator, in proton-proton collisions at center of mass energies of 7 TeV and 8 TeV [4,5]. The discovery of the Higgs boson, with a mass of 125 GeV, is a remarkable experimental verification of the SM, and contributes to rank the SM as one of the most successful theories developed to understand the fundamental components of matter and their interactions.

There are still important open questions that need an answer to get farther ahead in the understanding of nature; two of them are to decipher the dark matter (DM) enigma [6] and to solve the hierarchy problem of the SM [7–9]. These two problems, with no natural explanation within the SM, serve as a ground evidence that the SM might just well be the low energy approximation of a more complete theory. Supersymmetry (SUSY) is one of the most experimentally scrutinized models given the natural answers it provides to the DM and hierarchy questions [10,11]. The extended spectrum of SM particles, with new superpartners, which differ in spin with their SM counterparts, solves the hierarchy problem by the cancellation of quadratic divergent radiative corrections to the Higgs mass. The R-parity conserving SUSY models, such as the Minimal Supersymmetric Extension of the Standard Model (MSSM) have a genuine candidate for DM as the lightest SUSY particle (LSP) that in most of the scenarios is the lightest neutralino, $\tilde{\chi}_1^0$, given that it is electrically neutral and stable [12].

Squarks and gluinos, the strongly interacting superpartners of quarks and gluons with expected high production cross sections, were the candidates with the highest discovery potential at the startup of the LHC accelerator. After accumulating proton-proton collision data with an integrated luminosity of about 25 fb^{-1} in run 1 (years 2011 and 2012), and 120 fb^{-1} so far in run 2 (years 2015 through 2018), by each ATLAS and CMS experiments [13], no evidence of SUSY has been found yet. The large statistical samples recorded have allowed stringent exclusion limits on the masses of SUSY particles, by interpreting under different theory scenarios the results of non-excess-data over SM. At the moment, gluinos have the highest exclusion limits, reaching the 2 TeV mass border, and the third generation squarks masses are being excluded at the range above 1 TeV [14,15].

The neutralinos ($\tilde{\chi}_i^0$, $i=1\dots4$) and charginos ($\tilde{\chi}_2^\pm$, $\tilde{\chi}_1^\pm$), the mass eigenstates of the superpartners of the SM gauge and Higgs bosons, usually called electroweakinos since they are produced through electroweak processes, are expected to be lighter than squarks and gluinos in many SUSY models. They are characterized by much smaller production cross sections at hadron colliders, and therefore, their mass exclusion limits are lower than for colored SUSY particles. Limits are especially difficult to establish in the parameter space where the mass of the LSP is slightly lower than its SUSY mother

particle (compressed spectra). The search for electroweakinos is peculiarly interesting because of the connection between SUSY and DM. In the MSSM with R-parity conservation, SUSY particles are produced in pairs, and among all chargino-neutralino production processes, the $\tilde{\chi}_1^\pm/\tilde{\chi}_2^0$ electroweakino production usually has the highest branching ratio to leptons in the final state. The main chargino decays into leptonic states are: $\tilde{\chi}_1^\pm \rightarrow \tilde{\ell}\nu \rightarrow \ell\nu\tilde{\chi}_1^0$, $\tilde{\chi}_1^\pm \rightarrow \tilde{\ell}\nu \rightarrow \ell\tilde{\chi}_1^0\nu$, $\tilde{\chi}_1^\pm \rightarrow W\tilde{\chi}_1^0 \rightarrow \ell\nu\tilde{\chi}_1^0$. For the case of the next-to-lightest neutralino, the main leptonic decays are: $\tilde{\chi}_2^0 \rightarrow \tilde{\ell}^\pm\ell^\mp \rightarrow \ell^\pm\ell^\mp\tilde{\chi}_1^0$, $\tilde{\chi}_2^0 \rightarrow Z\tilde{\chi}_1^0 \rightarrow \ell^\pm\ell^\mp\tilde{\chi}_1^0$, $\tilde{\chi}_2^0 \rightarrow H\tilde{\chi}_1^0 \rightarrow WW\tilde{\chi}_1^0 \rightarrow \ell^\pm\ell^\mp\nu\bar{\nu}\tilde{\chi}_1^0$, $\tilde{\chi}_2^0 \rightarrow H\tilde{\chi}_1^0 \rightarrow ZZ\tilde{\chi}_1^0 \rightarrow \ell^\pm\ell^\mp\ell^\pm\ell^\mp\tilde{\chi}_1^0$.

The multilepton production, with no hadronic activity and with missing transverse momentum in the final decay chain of electroweakinos has been the main motivation of previous searches at the LHC (neutrinos and neutralinos are expected to escape without any detection, producing the missing momentum). Different studies have concentrated in a specific number of leptons in the final state, jet vetoes and dilepton electric charge [16, 17]. Nonetheless, the exclusion limits established by all these studies have limited sensitivity in the compressed mass region, where the mass splitting between $\tilde{\chi}_1^\pm/\tilde{\chi}_2^0$ and $\tilde{\chi}_1^0$ is very small ($\Delta m = m_{\tilde{\chi}_1^\pm} - m_{\tilde{\chi}_1^0} < 50$ GeV). For these low mass splittings the leptonic final states of $\tilde{\chi}_1^\pm/\tilde{\chi}_2^0$ are favoured through slepton decays with slepton masses in between $\tilde{\chi}_1^\pm/\tilde{\chi}_2^0$ and $\tilde{\chi}_1^0$. Usually, two study scenarios are considered, 1) stau dominated scheme (electroweakinos decaying preferentially to taus); 2) slepton democratic scheme (electroweakinos with equal probability for slepton decays). Note that in the compressed scenario, giving that the SUSY mother and daughter particles have a very small mass difference, the kinetic energy of the final products is limited; the final state leptons would in general have a soft p_T spectrum. Trigger and identification of leptons could reduce the efficiency to recognize them. In this case, the quest for SUSY could benefit from new tools.

Vector boson fusion (VBF) is the second main production process of Higgs bosons in proton-proton collisions at the LHC. Even though VBF Higgs production cross sections are about one order of magnitude lower than gluon-gluon fusion [18], VBF tagging technique has been effectively used to reduce backgrounds and allow Higgs properties to be measured with good precision in many Higgs physics analysis at the LHC [19]. The high luminosity samples provided by the LHC increase the feasibility of using VBF tagging in other physics studies. The VBF process consists on the radiation of vector bosons by incoming partons from each colliding proton. The partons remain intact after emission of the W bosons and then hadronize, keeping most of their initial energy and traveling very close to their initial trajectory. The radiated vector bosons fuse together through an electroweak process generating new particles, such as the Higgs or electroweakinos. The interesting feature of VBF is the existence of two forward very energetic jets that serve as a tagging scheme to reduce backgrounds. There is already one VBF SUSY search performed by the CMS experiment in run 1 of the LHC [20], which focused on the search of electroweakinos at $\sqrt{s} = 8$ TeV with 2 leptons in the final state. This study demonstrated the power of the VBF topology to investigate the SUSY compressed spectra scenarios increasing the sensitivity in this difficult SUSY parameter space. One critical aspect on the exploration of electroweakinos in the compressed SUSY space is the soft nature of the leptons in the final state. Making it difficult to identify them and reducing the final statistical sample when few leptons are required.

We present in this document the search of electroweakinos produced via VBF with one soft muon in the final state and missing transverse momentum. The SUSY topology is enhanced from backgrounds by the identification of the two VBF jets that provide the largest dijet invariant mass of

the event and the highest dijet pseudorapidity separation. We concentrate on the 2016 data sample recorded by the CMS experiment, which corresponds to a total integrated luminosity of 35.9 fb^{-1} of proton-proton collisions at center of mass energy of 13 TeV. We define control regions to contrast Monte Carlo (MC) simulations to data, correct any MC mismodeling and predict the event yields in the signal region defined by our event selection. The optimization of the event selections is motivated by MSSM electroweakinos produced with very low mass splitting with respect to the LSP mass and within an slepton democratic scenario. Our results are compared to the other single lepton analysis (electron, tau and when the final state leptons would not have enough energy to be reconstructed by the detector). Since no excess is observed over expected SM backgrounds we determine 95% CL exclusion limits in the SUSY parameter space spanned by the mass splitting Δm and the chargino mass.

The analysis presented in this document, combined with the other lepton channels with same VBF topology (identified in the CMS experiment as SUS-17-007), has been recently approved by the Analysis Review Committee and it is going right now through the Collaboration Wide review for final publication.

The author of this thesis participated in the upgrade for the startup of run 2 of one of the CMS muon subdetectors (RPC's), and in the CMS data collection and certification of the muon data quality. The author is a member of the CMS analysis group focused on SUSY searches through VBF and contributed to the group with the analysis of the VBF SUSY signal with one muon in the final state.

The present document is organized as follows: Chapters 1 and 2 are devoted to a brief review of the main aspects of the Standard model and supersymmetry, respectively. Chapter 3 describes the LHC accelerator and its operation. Chapter 4 provides a description of the CMS experiment and its subdetectors. Chapter 5 explains the physics object reconstruction and the main physics variables needed for our study. Chapter 6 reports all the analysis details, and Chapter 7 provides the results obtained.

Chapter 1

The Standard Model

The Standard Model (SM) of particle physics is a relativistic quantum field theory that summarizes our understanding of fundamental particles and their interactions, except gravity. The SM has been built upon experimental results. Nevertheless, it has made several predictions that have been later confirmed by experiments, making it a very successful model, so far. Despite of all its success, it cannot be considered a complete theory, because it does not include gravity nor dark matter, and in addition, it has several theoretical problems that suggest it is only a good approximation for low energies [21].

In the SM, particles are either bosons or fermions, the first ones with integer and the second with half-integer spin. Leptons ($e, \mu, \tau, \nu_e, \nu_\mu, \nu_\tau$) and quarks (u, d, t, b, c, s) are classified as fermions and are considered to be the basic building blocks of the Universe. In the SM, the fundamental interactions among particles are carried out through the exchange of gauge bosons. For every interaction there is an associated field, and for every field there is an associated quanta of the field: the gauge bosons. The photon is the gauge boson or carrier of the electromagnetic interaction, the vector bosons $W^\pm$ and Z^0 are the carriers of the weak interaction and the gluons are the carriers of the strong interaction (all of them with spin 1). In the case of the existence of a quantum field theory for gravity, the corresponding boson would be the graviton with spin 2. Table 1.1 summarizes the four fundamental interactions with their corresponding mediators, relative strength, and range. The

Table 1.1: Mediators, their spin, relative strength (with respect to the Strong interaction) and range for the fundamental interactions presented in nature. Gravity is not included in the SM because there is not a quantum field theory for it yet. So the gravitons are just hypothetical particles.

Interaction	Mediators	Mediator spin	Relative Strength	Range (m)
Strong	Gluons (g)	1	1	10^{-15}
Electromagnetic	Photon (γ)	1	10^{-3}	∞
Weak	$W^\pm, Z^0$	1	10^{-14}	10^{-18}
Gravitational	Gravitons?	2	10^{-43}	∞

first particle of the SM that was discovered is the electron (Cavendish Laboratory, 1897 [22]). The electron (e) and the electronic neutrino (ν_e) constitute the first generation of leptons. The second generation is conformed by the muon, μ , and the muon neutrino (ν_μ), whereas the third generation is made by the tau (τ) and the taonic neutrino (ν_τ). All the leptons have spin 1/2 and do not interact

strongly. In particular the neutral leptons (ν_e, ν_μ, ν_τ) interact only via the weak and force. Like leptons, quarks are also classified in three generations, but they interact strongly and consequently carry color charges. Figure 1.1 summarizes the building blocks of the Standard Model: fermions (leptons, quarks) and bosons with some of their basic properties (mass, electric charge and spin). All the ordinary matter in nature is composed of particles from the first generation: atomic nuclei are composed of quarks up and down and electrons are surrounding them as an electron cloud. The

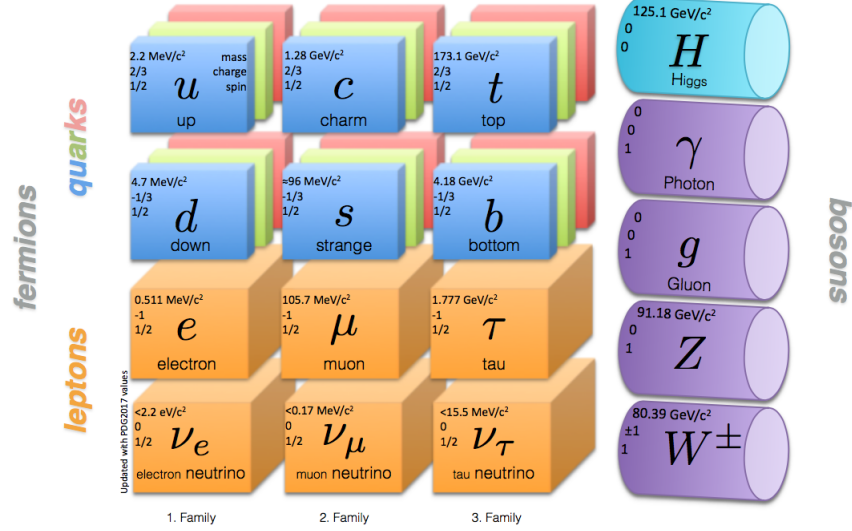


Figure 1.1: Particles of the SM, fermions are displayed in the three columns on the left and bosons in remaining column on the right. The properties of electric charge, spin and mass (or mass limit for neutrinos) are given in the corner of the boxes [23].

standard model also predicted the existence of the Higgs Boson. The Higgs boson is the quantum of the postulated Higgs field. The Higgs mechanism explains how particles acquire mass through their interaction with the Higgs field. Furthermore, it explains electroweak symmetry breaking (why the photon and gluons are massless while the $W^\pm$ and Z^0 are very heavy) [1–3]. On July 4th 2012, The ATLAS and CMS experiments of the Large Hadron Collider (LHC) announced a discovery of a scalar boson with a mass around $125 \text{ GeV}/c^2$ [4, 5], which was suspected to be the Higgs boson. Today, after many properties of this new boson have been measured by the LHC experiments, there is compelling evidence that confirms that the new boson discovered corresponds to the SM Higgs boson [24, 25].

1.1 Standard Model Symmetry Group

To build the Lagrangian of the Standard Model, some symmetries are imposed over the system. These symmetries are based on basic properties observed in nature. It means that every physical theory used to explain our universe has to satisfy the following properties:

- It has to be invariant under space time displacements
- It has to be invariant under rotations
- It has to be invariant under change of the inertial reference frame (boosts).

The last two cases form a group of symmetries called the Lorentz group. And the entire group of symmetries is known as the Poincare group. It means that the Lagrangian of the Standard Model has to be invariant under transformations of the Poincare group of symmetries. Other important feature of the SM Lagrangian is that it has to be renormalizable, this in order to avoid divergences when perturbative calculations are performed. Additionally, some local symmetries are required ($SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$, for the description of the strong, weak and electromagnetic interactions.

1.2 Gauge invariance and Quantum electrodynamics

One of the fundamental principles of the standard model is gauge invariance. To give an example, if we depart from the Dirac Lagrangian density for a spin 1/2 field [26]:

$$\mathcal{L}_0 = \bar{\psi}(x) i \not{\partial} \psi(x) - m \bar{\psi}(x) \psi(x) \quad (1.1)$$

In the previous equation we are using the Feynman usual notation [27]: $\not{\partial} := \gamma^\mu a_\mu$ so:

$$\not{\partial} \stackrel{\text{def}}{=} \gamma^\mu \partial_\mu \quad (1.2)$$

By applying a global phase transformation, the Lagrangian (and the resulting Dirac equation) remains invariant:

$$\psi(x) \longrightarrow \psi'(x) = e^{-iQ\varepsilon} \psi(x) \quad (1.3)$$

where $Q\varepsilon$ is an arbitrary real constant. But, what happens if the phase factor $Q\varepsilon$ depends of space and time?, in other words, what happens if we have a local phase transformation?

$$\psi(x) \longrightarrow \psi'(x) = e^{-iQ\varepsilon(x)} \psi(x) \quad (1.4)$$

We can see that the mass term in Equation 1.1 remains unchanged, but we have a problem with the kinetic term:

$$\partial_\mu \psi(x) \longrightarrow e^{-iQ\varepsilon(x)} (\partial_\mu - iQ\partial_\mu \varepsilon(x)) \psi(x) . \quad (1.5)$$

So, we have an extra term due to the dependence of ε on space-time ($\partial_\mu \varepsilon(x)$). This implies that if we want an invariant Lagrangian under transformations such as 1.3, the phase has to be the same in all the space-time. Since the Lagrangian from Equation 1.1 is not invariant under local phase transformations, we add an additional term to try to restore the invariance. So we introduce a spin-1 vector field A_μ with the idea to cancel the additional term of Equation 1.5. To achieve this, we replace the ordinary derivative ∂_μ by a covariant derivative D_μ :

$$D_\mu \psi(x) = (\partial_\mu + iQA_\mu(x)) \psi(x) , \quad (1.6)$$

with the transformation

$$A_\mu(x) \longrightarrow A'_\mu(x) = A_\mu(x) + \partial_\mu \varepsilon(x) . \quad (1.7)$$

We can check that the new covariant derivative takes the form:

$$D_\mu \psi(x) \longrightarrow (D_\mu \psi)'(x) = e^{-iQ\varepsilon(x)} D_\mu \psi(x) , \quad (1.8)$$

Now the modified Lagrangian is given by:

$$\mathcal{L} = \bar{\psi}(x) (i \not{D} - m) \psi(x) = \mathcal{L}_0 - QA_\mu(x) \bar{\psi}(x) \gamma^\mu \psi(x) \quad (1.9)$$

This new Lagrangian is invariant under local phase transformation (gauge invariant). We can see that from the gauge invariant requirement we introduced a new vector field A_μ that couples to the fermion field ψ (last term of Equation 1.9). However, in this Lagrangian, a kinetic term is missing for the field A_μ . So, including the kinetic term, we get the full Lagrangian given by:

$$\mathcal{L} = \bar{\psi} (i \not{D} - m) \psi - \frac{1}{4} F_{\mu\nu} F^{\mu\nu} \quad (1.10)$$

This Lagrangian remains gauge invariant because the field strength tensor $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ is also gauge invariant. Moreover, replacing Q with $-e$ for the electron field, we arrive to the quantum electrodynamics (QED) Lagrangian. As a summary, we start from the Lagrangian for a free spin 1/2 field and requiring gauge invariance, we arrive to the QED Lagrangian. The main conclusion is that the gauge invariance requirement generates the interaction between the fields. Yang and Mills used the same strategy in 1954 [28], insisting that a global invariance holds locally to the SU(2) group. Later, that idea was applied to the color SU(3) obtaining the gauge field theory for the strong interaction, quantum chromodynamics (QCD). All the fundamental interactions in the Standard Model are generated in this way.

1.3 Quantum chromodynamics

The gauge field theory that models the strong interactions of colored quarks and gluons is called quantum chromodynamics (QCD). From the SM $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$, QCD is the $SU(3)_C$ component. The observation of bound states such as $\Delta_{(uuu)}^{++}$ in 1951 [29] needed an explanation to avoid the violation of Pauli's exclusion principle [30]. It was necessary to introduce a new quantum number associated with the group SU(3): color [31, 32]. QCD has the following postulates:

- **Confinement:** color charged particles (like quarks and gluons) cannot be isolated, implying that they can not be observed isolated in nature [33]. When there is an attempt to separate them, they immediately clump together to form colorless bound states (hadrons) that can be mesons (one quark, one antiquark) or baryons (three quarks) in a process called hadronization. In other words, a shower of color neutral particles (called jet) is observed instead of a single quark. As a result of this, the range of the strong force is confined despite the fact that gluons are massless.
- **Asymptotic freedom:** the interaction of quarks becomes asymptotically weaker as the energy scale increases and the corresponding length scale decreases [34, 35]. This implies that perturbative corrections are allowed at high energies. Oppositely, the interaction becomes strong at low energies, leading to the creation of composite hadrons due to confinement of quarks and gluons, the perturbative calculations are not valid in this case.

QCD is based on the Special Unitary group in 3 (complex) dimensions, whose elements are the set of unitary 3×3 matrices with determinant one [36]. A complex $N \times N$ matrix has $2N^2$ degrees of freedom, but the unitarity requirement ¹ restricts the number of degrees of freedom to N^2 . This means that there are 9 linearly independent unitary complex matrices, but one has determinant -1 , so finally we have 8 different generators. These matrices can operate in a set of 3-vectors or triplets (which represent quarks in the color space) and also on themselves (which represent gauge transformations carried out by gluons) [38]. So, quarks are triplets in the fundamental representation of the SU(3) color group, while gluons are octets and transform under the adjoint representation of the SU(3) color group [23]. The QCD Lagrangian density is given by

$$\mathcal{L} = \bar{\psi}_q^i (i\gamma^\mu) (D_\mu)_{ij} \psi_q^j - m_q \bar{\psi}_q^i \psi_{qi} - \frac{1}{4} F_{\mu\nu}^a F^{a\mu\nu} , \quad (1.11)$$

The different parts of the equation have the following meaning:

- ψ_q^i represents a quark field spinor for a quark of flavor q and mass m_q with a color index i that runs from 1 to 3.

$$\psi_q = \begin{pmatrix} \psi_{qR} \\ \psi_{qG} \\ \psi_{qB} \end{pmatrix} \quad (1.12)$$

- γ^μ is a Dirac matrix with μ being a Lorentz vector index:

$$\begin{aligned} \gamma^0 &= \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & -1 & 0 \\ 0 & 0 & 0 & -1 \end{pmatrix}, & \gamma^1 &= \begin{pmatrix} 0 & 0 & 0 & 1 \\ 0 & 0 & 1 & 0 \\ 0 & -1 & 0 & 0 \\ -1 & 0 & 0 & 0 \end{pmatrix} \\ \gamma^2 &= \begin{pmatrix} 0 & 0 & 0 & -i \\ 0 & 0 & i & 0 \\ 0 & i & 0 & 0 \\ -i & 0 & 0 & 0 \end{pmatrix}, & \gamma^3 &= \begin{pmatrix} 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & -1 \\ -1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \end{pmatrix}. \end{aligned}$$

- D_μ is the covariant derivative in QCD defined by:

$$(D_\mu)_{ij} = \delta_{ij} \partial_\mu - i g_s t_{ij}^a A_\mu^a , \quad (1.13)$$

where:

- g_s is the strong coupling (related to α_s by $g_s^2 = 4\pi\alpha_s$)
- A_μ^a is the gluon field with color index a

¹A complex square matrix U is unitary if its conjugate transpose U^* is also its inverse: $U^*U = UU^* = I$. Where I is the identity matrix [37].

- t_{ij}^a is proportional to the hermitean and traceless Gell-Mann matrices of SU(3) (the generators) $t_{ij}^a = \frac{1}{2}\lambda_{ij}^a$ given by:

$$\begin{aligned}\lambda^1 &= \begin{pmatrix} 0 & 1 & 0 \\ 1 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, & \lambda^2 &= \begin{pmatrix} 0 & -i & 0 \\ i & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, & \lambda^3 &= \begin{pmatrix} 1 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & 0 \end{pmatrix}, & \lambda^4 &= \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 0 \\ 1 & 0 & 0 \end{pmatrix} \\ \lambda^5 &= \begin{pmatrix} 0 & 0 & -i \\ 0 & 0 & 0 \\ i & 0 & 0 \end{pmatrix}, & \lambda^6 &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix}, & \lambda^7 &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & -i \\ 0 & i & 0 \end{pmatrix}, & \lambda^8 &= \begin{pmatrix} \frac{1}{\sqrt{3}} & 0 & 0 \\ 0 & \frac{1}{\sqrt{3}} & 0 \\ 0 & 0 & -\frac{2}{\sqrt{3}} \end{pmatrix}\end{aligned}$$

These matrices represent the fact that the interaction between a quark and a gluon rotates the color of the quark in the SU(3) space [23].

- $F_{\mu\nu}^a$ is the gluon field strength tensor for a gluon with color index a (that means, $a \in [1, \dots, 8]$) defined by:

$$F_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a + g_s f^{abc} A_\mu^b A_\nu^c \quad (1.14)$$

The factors f^{abc} are called the structure constants of SU(3) given by [38]:

$$f_{123} = 1 \quad (1.15)$$

$$f_{147} = f_{246} = f_{257} = f_{345} = \frac{1}{2} \quad (1.16)$$

$$f_{156} = f_{367} = -\frac{1}{2} \quad (1.17)$$

$$f_{458} = f_{678} = \frac{\sqrt{3}}{2} \quad (1.18)$$

Antisymmetric in all indices

All other $f_{abc} = 0$

These factors are related to the generators of the fundamental representation t_{ij}^a through the commutator relations

$$t^a t^b - t^b t^a = [t^a, t^b] = i f^{abc} t^c, \quad (1.19)$$

There is the possibility to express the color space in two different ways, in the fundamental representation with the t_{ij}^a matrices, or with the structure constants f^{abc} [38].

It is important to notice that the fundamental parameters of QCD are the quark masses m_q and the coupling g_s (or $\alpha_s = \frac{g_s^2}{4\pi}$). Other important remark is that hadrons are color neutral combinations of quarks, antiquarks and gluons. That means, they do not carry color charge since quarks and gluons can not be observed as free particles due to confinement [23].

1.4 The Higgs mechanism

To illustrate the Higgs mechanism, we consider the case of a complex scalar field ϕ coupled to an electromagnetic field A_μ and to itself. The Lagrangian who describes this case is given by

$$\mathcal{L} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + D_\mu\phi^*D^\mu\phi - V(\phi) \quad (1.20)$$

where

$$F_{\mu\nu} = \partial_\nu A_\mu - \partial_\mu A_\nu. \quad (1.21)$$

D_μ is the covariant derivative

$$D_\mu = \partial_\mu - ieA_\mu \quad (1.22)$$

The mass term for the electromagnetic field ($\propto m^2 A_\mu A^\mu$), which in this case represents the photon mass, is forbidden in the Lagrangian. Since this Lagrangian is required to be invariant under local gauge transformations like these

$$\phi(x) \rightarrow e^{i\alpha(x)}\phi(x) \quad A_\mu(x) \rightarrow A_\mu(x) - \frac{1}{e}\partial_\mu\alpha(x) \quad (1.23)$$

for any α and x . And this photon mass term violates the local gauge invariance. The most general renormalizable potential $V(\phi)$ [39] is given by

$$V(\phi) = \mu^2\phi^*\phi + \lambda(\phi^*\phi)^2. \quad (1.24)$$

An additional requirement for this potential is that λ has to be positive in order to have a global minimum [40]. We need the potential to remain finite as ϕ goes to $\pm\infty$, that means, the potential has to be bounded. Now, for the μ^2 term we have two options:

- $\mu^2 > 0$, in this case the minimum of the potential is obtained for $\langle 0|\phi|0\rangle \equiv \phi_0 = 0$. In other words, the lowest energy state is the vacuum, as shown in the left side of Figure 1.2. In this case, Equation 1.20 represents the QED Lagrangian for a scalar charged particle with mass μ and self-interactions depicted by the term ϕ^4 .
- $\mu^2 < 0$, for this case we obtain the minimum for the potential by equalizing the derivative of the potential to zero, which yields

$$\langle \phi \rangle_0 \equiv \langle 0|\phi|0\rangle = \sqrt{-\frac{\mu^2}{2\lambda}} \equiv \frac{v}{\sqrt{2}} \quad (1.25)$$

the minimum is not anymore at $\phi = 0$, as in the previous case. The term $\langle 0|\phi|0\rangle$ is known as the vacuum expectation value (VEV) for the field ϕ . This VEV is chosen conventionally in the direction of the real part of the field ϕ . However, this choice is arbitrary, another direction can be chosen for the VEV. The right side of Figure 1.2 shows the location of the minimum in the case where the scalar field ϕ is real. As a conclusion, the requirement of $\mu^2 < 0$ is breaking the symmetry introducing a vacuum expectation value different from zero (which is chosen arbitrarily but can be in any direction of the complex scalar field ϕ).

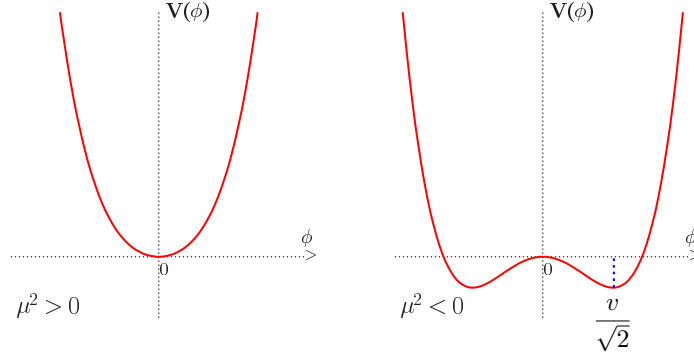


Figure 1.2: The potential V of the scalar field ϕ in the case $\mu^2 > 0$ (left) and $\mu^2 < 0$ (right). Extracted from [41].

The next step is to expand the Lagrangian around the vacuum state $\langle \phi \rangle_0$, the minima (the calculations for the expansion are detailed in Reference [42])

$$\phi(x) = \frac{1}{\sqrt{2}} (v + \eta(x) + i\xi(x)) \quad (1.26)$$

to perform the expansion it is assumed that

$$\langle 0 | \eta | 0 \rangle = \langle 0 | \xi | 0 \rangle = 0. \quad (1.27)$$

With the definition from 1.26, the product of the covariant derivatives is given by

$$\begin{aligned} (D_\mu \phi)^* (D^\mu \phi) &= \frac{1}{2} \partial_\mu \eta \partial^\mu \eta + \frac{1}{2} \partial_\mu \xi \partial^\mu \xi + \frac{e^2 v^2}{2} A_\mu A^\mu + \frac{e^2}{2} A_\mu A^\mu (2v\eta + \eta^2 + \xi^2) - ev A^\mu \partial_\mu \xi \\ &+ e A^\mu \xi \partial_\mu \eta - e A^\mu \eta \partial_\mu \xi. \end{aligned}$$

And the potential becomes:

$$\begin{aligned} V(\phi) &= \mu^2 \phi^* \phi + \lambda (\phi^* \phi)^2 = \frac{\mu^2}{2} (v^2 + 2v\eta + \eta^2 + \xi^2) \\ &+ \frac{\lambda}{4} (v^4 + \eta^4 + \xi^4 + 4v^3\eta + 6v^2\eta^2 + 2v^2\xi^2 + 4v\eta^3 + 4v\eta\xi^2 + 2\eta^2\xi^2). \end{aligned}$$

Using the relation $\mu^2 = -v^2\lambda$ the potential now is given by:

$$V(\phi) = -\frac{v^4\lambda}{4} + v^2\lambda\eta^2 + \frac{\lambda}{4}\eta^4 + \frac{\lambda}{4}\xi^4 + v\lambda\eta^3 + v\lambda\eta\xi^2 + \frac{\lambda}{2}\eta^2\xi^2.$$

Finally, our departing Lagrangian from Equation 1.20 is transformed into:

$$\begin{aligned}\mathcal{L} = & \frac{v^4\lambda}{4} - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{e^2v^2}{2}A_\mu A^\mu + \frac{1}{2}\partial_\mu\eta\partial^\mu\eta - \frac{2v^2\lambda}{2}\eta^2 + \frac{1}{2}\partial_\mu\xi\partial^\mu\xi - evA^\mu\partial_\mu\xi \\ & + \frac{e^2}{2}A_\mu A^\mu(2v\eta + \eta^2 + \xi^2) + eA^\mu\xi\partial_\mu\eta - eA^\mu\eta\partial_\mu\xi - \frac{\lambda}{4}(\eta^2 + \xi^2)^2 - v\lambda\eta(\eta^2 + \xi^2).\end{aligned}\quad (1.28)$$

From the previous Lagrangian, we identify the quadratic terms in the fields as the mass terms. Then we have a term for the mass of the photon

$$\frac{1}{2}M_A^2 A_\mu A^\mu = \frac{e^2v^2}{2} A_\mu A^\mu \quad (1.29)$$

that yields

$$M_A = ev = -\frac{e\mu^2}{\lambda}. \quad (1.30)$$

In a similar way

$$\frac{1}{2}M_\eta^2\eta^2 = \frac{2v^2\lambda}{2}\eta^2 \quad (1.31)$$

yields

$$M_\eta^2 = 2v^2\lambda = -2\mu^2 \quad (1.32)$$

Apparently, for the field ξ we have a massless particle since there is no quadratic term in the fields for ξ . The appearance of this massless field allows us to formulate the Goldstone theorem that states that when a global symmetry is spontaneously broken (as in the previous case), one or several bosons without mass (spin 0), called Goldstone bosons, are obtained. The number of Goldstone bosons depend on the number of generators broken by the spontaneous symmetry breaking. However, in the previous case, we have a problem with the degrees of freedom in the Lagrangian. At the beginning, in Equation 1.20, we have four degrees of freedom, two for the electromagnetic field (without mass) A^μ and two for the complex scalar field ϕ . However, with the spontaneous symmetry breaking, we have now five degrees of freedom, three for the massive A^μ photon, and one for η and one for ξ . This tells us that there is an additional field that has no physical meaning. Additionally, the Lagrangian has a quadratic term (that must be eliminated) that includes the fields A^μ and ξ : $evA^\mu\partial_\mu\xi$. To achieve this, we use the fact that at first order:

$$\phi(x) = \frac{1}{\sqrt{2}}(v + \eta + i\xi) \equiv \frac{1}{\sqrt{2}}[v + \eta'(x)]e^{i\zeta(x)/v} \quad (1.33)$$

With real $\eta'(x)$ and $\zeta(x)$. And additionally we use the gauge transformation:

$$\phi(x) \rightarrow e^{-i\zeta(x)/v}\phi(x), \quad A_\mu \rightarrow A_\mu - \frac{1}{ev}\partial_\mu\zeta(x) \quad (1.34)$$

that yields the following Lagrangian

$$\mathcal{L} = \frac{1}{2}\partial_\mu\eta\partial^\mu\eta - \frac{2v^2\lambda}{2}\eta^2 - \frac{F_{\mu\nu}F^{\mu\nu}}{4} + \frac{e^2v^2}{2}A_\mu A^\mu + \frac{e^2}{2}A_\mu A^\mu(2v\eta + \eta^2) + \frac{v^4\lambda}{4} - \lambda v\eta^3 - \frac{\lambda}{4}\eta^4. \quad (1.35)$$

With the previous choice, all the terms with ξ disappear from the Lagrangian. That is, this gauge transformation (called unitary gauge) has led us to obtain only physical states in the Lagrangian.

The photon, initially with two degrees of freedom and without mass, has absorbed what seems to be a Goldstone boson (the field ξ) and has become a massive field with three degrees of freedom. As a conclusion, this allows us to elucidate how the Higgs mechanism works. Initially, we have a spontaneous symmetry breaking in which Goldstone bosons are generated. These bosons through a unitary gauge transformation are absorbed (disappearing from the Lagrangian) resulting in the appearance of massive bosonic fields.

1.5 Electroweak sector

The electroweak interaction describes two of the four known fundamental interactions of nature: electromagnetism and the weak interaction. Even though these forces are different at low energies, in reality they are different aspects of the same force: the electroweak force [43–45]. In principle, above the unification energy (~ 246 GeV), the electromagnetic force and weak force would merge into a combined electroweak force. The electroweak interactions were experimentally proven first by the Gargamelle collaboration in 1973 with the discovery of neutral currents in neutrino scattering [46–48], and second, with the discovery of the W and Z gauge bosons at the Super Proton Synchrotron in proton-antiproton collisions [49–52]. The electroweak Lagrangian density can be written as a sum of several contributions. One of these contribution is $\mathcal{L}_{\text{gauge}}$, that describes the interaction between three $SU(2)_L$ gauge bosons, W_μ^a , $a = 1, 2, 3$, and one $U(1)_Y$ gauge boson, B_μ and can be written as:

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4}W_{\mu\nu}^a W_a^{\mu\nu} - \frac{1}{4}B_{\mu\nu}B^{\mu\nu}, \quad (1.36)$$

where,

$$\begin{aligned} W_{\mu\nu}^a &= \partial_\nu W_\mu^a - \partial_\mu W_\nu^a + g\epsilon^{abc}W_\mu^b W_\nu^c, \\ B_{\mu\nu} &= \partial_\nu B_\mu - \partial_\mu B_\nu \end{aligned} \quad (1.37)$$

In addition, the coupling constant for $SU(2)_L$ is g and for $U(1)_Y$ is g' . Also, a complex scalar $SU(2)$ doublet, Φ is coupled to the gauge fields:

$$\Phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix}. \quad (1.38)$$

The justification of the introduction of the complex scalar doublet is coming from the Higgs mechanism explained in the previous section. Similarly as Equation 1.24, the potential for the complex scalar doublet of Equation 1.38 is:

$$V(\Phi) = \mu^2 |\Phi^\dagger \Phi| + \lambda \left(|\Phi^\dagger \Phi| \right)^2, \quad (1.39)$$

with $\lambda > 0$. In this way, the scalar contribution to the Lagrangian density is:

$$\mathcal{L}_{\text{scalar}} = (D^\mu \Phi)^\dagger (D_\mu \Phi) - V(\Phi), \quad (1.40)$$

where the covariant derivative that leaves this Lagrangian invariant under gauge transformations is:

$$D_\mu = \partial_\mu - ig\frac{\tau_a}{2}W_\mu^a - ig'\frac{1}{2}B_\mu Y \quad (1.41)$$

τ_a represents the Pauli matrices given by:

$$\tau_1 = \tau_x = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix} \quad (1.42)$$

$$\tau_2 = \tau_y = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix} \quad (1.43)$$

$$\tau_3 = \tau_z = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}. \quad (1.44)$$

As in the previous Section, the state of minimum energy for $\mu^2 < 0$ is not at $\phi^0 = 0$ (as shown in Figure 1.3) and the scalar field develops a vacuum expectation value. Since the potential depends only on the combination $\Phi^\dagger \Phi$, the direction of the minimum in $SU(2)$ space is not determined. So, we can arbitrarily choose

$$\langle \Phi \rangle_0 \equiv \langle 0 | \Phi | 0 \rangle = \begin{pmatrix} 0 \\ \frac{v}{\sqrt{2}} \end{pmatrix} \quad \text{with } v = \left(-\frac{\mu^2}{\lambda} \right)^{1/2} \quad (1.45)$$

where we assign hypercharge $Y = 1$ to Φ . Continuing with the field Φ , we write it in terms of four

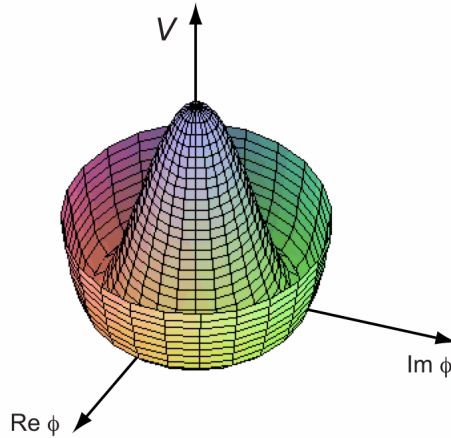


Figure 1.3: The potential ϕ with an unstable state at $\phi = 0$ and minima around a circle. [53].

fields $\theta_{1,2,3}(x)$ and $H(x)$ at first order:

$$\Phi(x) = \begin{pmatrix} \theta_2 + i\theta_1 \\ \frac{1}{\sqrt{2}}(v + H) - i\theta_3 \end{pmatrix} = e^{i\theta_a(x)\tau^a(x)/v} \begin{pmatrix} 0 \\ \frac{1}{\sqrt{2}}(v + H(x)) \end{pmatrix} \quad (1.46)$$

And now we perform a gauge transformation on this field, with the idea of picking the gauge for which only the physical particles are left in the Lagrangian (unitary gauge selection)

$$\Phi(x) \rightarrow e^{-i\theta_a(x)\tau^a(x)} \Phi(x) = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + H(x) \end{pmatrix} \quad (1.47)$$

Now, writing the term $|D_\mu \Phi|^2$ (the first term of the Lagrangian, Equation 1.40), we get:

$$\begin{aligned}
|D_\mu \Phi|^2 &= \left| \left(\partial_\mu - ig \frac{\tau_a}{2} W_\mu^a - ig' \frac{1}{2} B_\mu \right) \Phi \right|^2 \\
&= \frac{1}{2} \left| \begin{pmatrix} \partial_\mu - \frac{i}{2}(gW_\mu^3 + g'B_\mu) & -\frac{ig}{2}(W_\mu^1 - iW_\mu^2) \\ -\frac{ig}{2}(W_\mu^1 + iW_\mu^2) & \partial_\mu + \frac{i}{2}(gW_\mu^3 - g'B_\mu) \end{pmatrix} \begin{pmatrix} 0 \\ v + H \end{pmatrix} \right|^2 \\
&= \frac{1}{2} (\partial_\mu H)^2 + \frac{1}{8} g^2 (v + H)^2 |W_\mu^1 + iW_\mu^2|^2 + \frac{1}{8} (v + H)^2 |gW_\mu^3 - g'B_\mu|^2
\end{aligned} \tag{1.48}$$

Based on Equation 1.48, we define the new fields $W_\mu^\pm$ and Z_μ (these fields are orthogonal to A_μ) as:

$$W^\pm = \frac{1}{\sqrt{2}} (W_\mu^1 \mp iW_\mu^2), \quad Z_\mu = \frac{gW_\mu^3 - g'B_\mu}{\sqrt{g^2 + g'^2}}, \quad A_\mu = \frac{gW_\mu^3 + g'B_\mu}{\sqrt{g^2 + g'^2}} \tag{1.49}$$

the terms which are bilinear in the fields $W^\pm, Z, A$ represent the masses:

$$M_W^2 W_\mu^+ W^{-\mu} + \frac{1}{2} M_Z^2 Z_\mu Z^\mu + \frac{1}{2} M_A^2 A_\mu A^\mu \tag{1.50}$$

So, the photon is still massless while the W and Z bosons have acquired masses

$$M_W = \frac{1}{2} v g, \quad M_Z = \frac{1}{2} v \sqrt{g^2 + g'^2}, \quad M_A = 0 \tag{1.51}$$

We obtain in this way, the desired spontaneously symmetry breaking $SU(2)_L \times U(1)_Y \rightarrow U(1)_Q$. In addition the three Goldstone bosons produced are absorbed by the $W^\pm$ and Z bosons to form their longitudinal components and to get their masses [41]. The photon, that is the generator of the $U(1)_Q$ remains massless because the symmetry of $U(1)_Q$ is still unbroken. An important remark is that there is no prediction over the $W^\pm$ and Z boson masses since g and g' are free parameters of the model. We can use the same scalar field Φ with hypercharge $Y=1$, and the isodoublet $\tilde{\Phi} = i\tau_2 \Phi^*$, which has hypercharge $Y=-1$ to generate the fermion masses. For every fermion generation we can add the $SU(2)_L \times U(1)_Y$ invariant Yukawa Lagrangian:

$$\mathcal{L}_F = -\lambda_e \bar{L} \Phi e_R - \lambda_d \bar{Q} \Phi d_R - \lambda_u \bar{Q} \tilde{\Phi} u_R + h.c. \tag{1.52}$$

and perform the same procedure followed for the vector bosons W and Z .

In case of the electron, we have

$$\begin{aligned}
\mathcal{L}_F &= -\frac{1}{\sqrt{2}} \lambda_e (\bar{\nu}_e, \bar{e}_L) \begin{pmatrix} 0 \\ v + H \end{pmatrix} e_R + \dots \\
&= -\frac{1}{\sqrt{2}} \lambda_e (v + H) \bar{e}_L e_R + \dots
\end{aligned} \tag{1.53}$$

Again, the terms that are bilinear in the fermion fields $\bar{f}_L f_R$ (in particular for the electron $\bar{e}_L e_R$) represent the fermion masses

$$m_e = \frac{\lambda_e v}{\sqrt{2}}, \quad m_u = \frac{\lambda_u v}{\sqrt{2}}, \quad m_d = \frac{\lambda_d v}{\sqrt{2}} \tag{1.54}$$

As a conclusion, the addition of the scalar doublet Φ generated the masses of fermions and vector bosons $W^\pm, Z$, at the same time, the $SU(2) \times U(1)$ gauge symmetry is now spontaneously broken. Nevertheless, the $SU(3)$ color symmetry and the electromagnetic $U(1)_Q$ symmetry stay unbroken [41].

1.6 Experimental Results of the Standard Model

All the particles predicted by the standard model have been discovered. The quark sector was completed with the discovery of the top quark in 1995 by the CDF and D0 experiments of the Tevatron collider [54,55]. The lepton sector got completed with the discovery of the tau neutrino by the DONUT collaboration of FERMILAB laboratory in 2000 [56]. After many searches in different particle accelerators, the last boson proposed by the SM was discovered in 2012 by the ATLAS and CMS collaboration of the LHC accelerator. It was initially announced as a discovery of a resonance with a mass of about 125 GeV [4, 5], with characteristics very similar to those of the Higgs boson proposed by the Standard Model. The LHC Run 1 pp collision data set consisted of an integrated luminosity of 5 fb^{-1} accumulated by each, ATLAS and CMS, at $\sqrt{s} = 7 \text{ TeV}$, and 20 fb^{-1} recorded at $\sqrt{s} = 8 \text{ TeV}$ by each experiment [57]. Subsequent studies performed with these two data samples about the properties of the scalar boson discovered, such as decay rates, production properties, couplings to SM particles, spin, parity, width and lifetime have convinced the scientific community that the bosonic resonance discovered was actually the SM Higgs boson. With this discovery, the bosonic content of the SM is complete. The full particle spectrum predicted by the SM being fully discovered provides solid ground evidence that the SM is the most successful theory to explain nature compositeness and the interaction of its fundamental components. Figure 1.4 shows the invariant mass distributions for the main discovery channels of the Higgs, $H \rightarrow \gamma\gamma$, and $H \rightarrow ZZ \rightarrow 4\ell$, with a well defined resonance observed on top of smooth backgrounds from other well understood SM processes. Table 1.2 shows the most important decay channels to study the Higgs boson properties in the LHC. The table shows the channels that present better resolution in the reconstruction of the Higgs mass, $H \rightarrow \gamma\gamma$ and $H \rightarrow ZZ \rightarrow 4\ell$. For these channels, all the final states can be accurately measured. For the channel $H \rightarrow W^+W^- \rightarrow \ell^+\nu_\ell\ell'^-\bar{\nu}_{\ell'}$, there is a reduction in mass resolution due to the presence of neutrinos in the final state. The channels $H \rightarrow b\bar{b}$ and $H \rightarrow \tau^+\tau^-$ have also a low mass resolution, given that the final states are more difficult to reconstruct, due to the existence of large backgrounds mimicking the signal [23]. An example of the accuracy of the measurements achieved by the LHC experiments with the run1 data is the higgs mass: 125.09 ± 0.21 (statistical) ± 0.11 (systematic). The fact that statistical uncertainties are still dominant compared to systematics, motivates the need to gather a data sample with higher integrated luminosity, something that is being achieved with the present LHC run 2 undergoing at the moment. Even though the Higgs studies have driven most of the attention, they are just

Decay channel	Mass resolution
$H \rightarrow \gamma\gamma$	1–2%
$H \rightarrow ZZ \rightarrow \ell^+\ell^-\ell'^+\ell'^-$	1–2%
$H \rightarrow W^+W^- \rightarrow \ell^+\nu_\ell\ell'^-\bar{\nu}_{\ell'}$	20%
$H \rightarrow b\bar{b}$	10%
$H \rightarrow \tau^+\tau^-$	15%

Table 1.2: The five main decay channels for the study of the SM Higgs boson at the LHC [23].

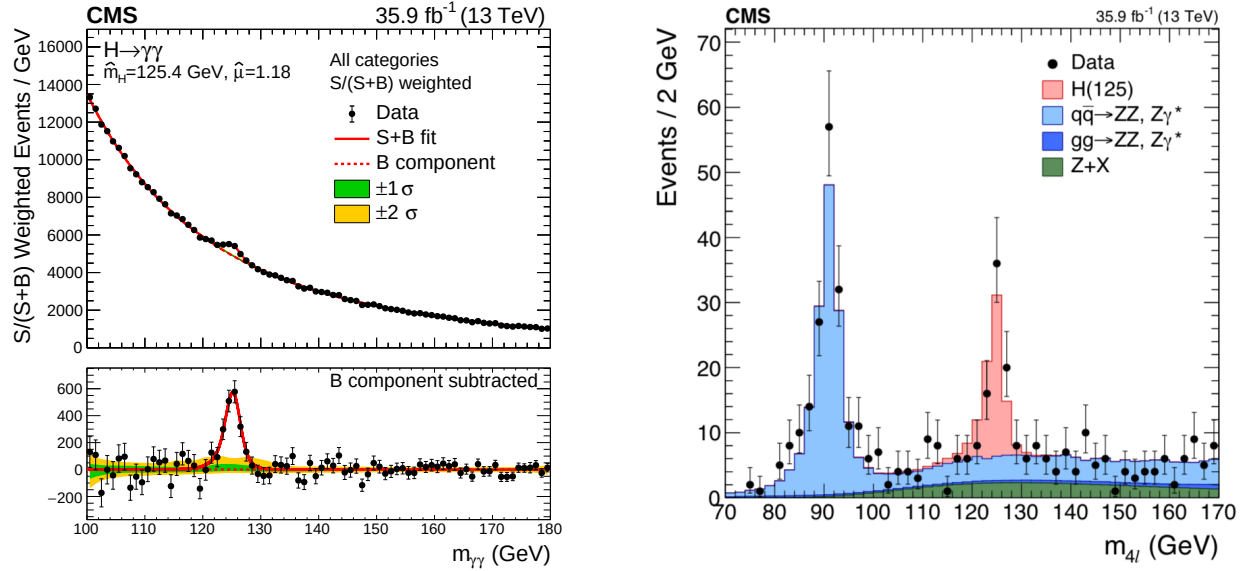


Figure 1.4: (Left) The invariant mass distribution of diphoton candidates, with each event weighted by the signal-to-background ratio in each event category, observed by CMS [25] at Run 2. In the lower panel are displayed the residuals of the data with respect to the fitted background. (Right) The m_{4l} distribution from CMS Run 2 data [24].

part of a huge effort to test experimentally the SM with high precision measurements. Figure 1.5 depicts a summary of production cross section measurements, performed by the CMS experiment, for different SM processes. The red and blue boxes show the measured values for the cross sections obtained during Run 1 of the LHC with energies of $\sqrt{s} = 7$ and 8 TeV respectively (the green boxes indicate measurements performed at $\sqrt{s} = 13$ TeV). In cases where there are no boxes but one line, it is because the indicated process has not been observed yet, in that case, 95% C.L exclusion limits are indicated. One of the cases that attracts attention is the production of the Higgs in association with a pair of top-antitop quarks. Reference [59], published recently, shows that the observation of $t\bar{t}H$ production is in excellent agreement with the Standard Model. This observation is the first confirmation of the tree-level coupling of the Higgs boson to top quarks. Figure 1.6 shows examples of tree-level Feynman diagrams for the production process of Higgs in association with of a pair of top-antitop quarks in proton-proton collisions. This is a strong indication that the Higgs boson has a crucial role in the large value of the top quark mass. It is the first time that this essential feature of the Standard Model has been verified experimentally with high significance. The study presented in this document requires the knowledge of few of SM processes, we concentrate on describing the main features of these processes.

1.6.1 W and Z bosons

The $W^\pm$ bosons are mediators of the neutrino emission or absorption. They are responsible for processes where flavor or lepton change is present. A strange quark transformed into an up quark or the muon decay to an electron and two neutrinos are examples of this kind of processes. The

charge carried by the W boson induces the emission or absorption of electrons or positrons, as can be seen in Figure 1.7. In contrast, the Z boson is not involved in the absorption/emission of neutrinos. The Z boson participates as a momentum transfer mediator in weak processes like the scattering of neutrinos with matter. An electron that suddenly increases its kinetic energy in the presence of a neutrino beam is an example of such kind of mediation. The increase is understood as the neutrino interaction with the electron mediated by a Z boson emission. Table 1.3 displays the average world measurements for mass, width, spin, and branching fractions of the W and Z boson. Table 1.3 also shows the cross sections measured by the CMS experiment at $\sqrt{s}=13$ TeV [60]. Note that the W^+ production cross section from proton proton collisions is $\sim 30\%$ higher than W^- . The reason for this difference is that protons have a higher “up” quark content and therefore the probability of generating W^+ bosons is higher. Figure 1.8 shows the typical diagrams for the production of W+jets

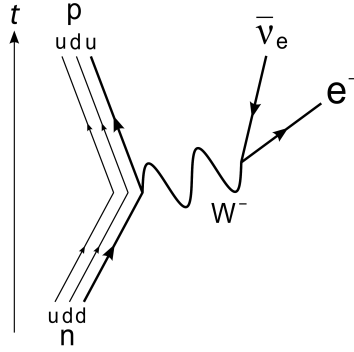


Figure 1.7: Feynman diagram for beta decay of a neutron into a proton, electron, and electron antineutrino via a W boson [61].

Table 1.3: Average world measurements for mass, width, spin, and branching fractions of the W boson and Z boson, obtained from the particle data group [13] . Cross sections values measured by the CMS experiment at $\sqrt{s}=13$ TeV [60].

Properties	W Boson	Z Boson
Mass (GeV)	80.379 ± 0.012	91.1876 ± 0.0021
Width (GeV)	2.085 ± 0.042	2.4952 ± 0.0023
Spin	1	1
Decay Modes (Branching Fractions)	$W^+ \rightarrow e^+ \nu$ (10.71 $\pm$ 0.16)% $W^+ \rightarrow \mu^+ \nu$ (10.63 $\pm$ 0.15)% $W^+ \rightarrow \tau^+ \nu$ (11.38 $\pm$ 0.21)% $W^+ \rightarrow \text{hadrons}$ (67.41 $\pm$ 0.25)%	$Z \rightarrow e^+ e^-$ (3.3632 $\pm$ 0.0042)% $Z \rightarrow \mu^+ \mu^-$ (3.3662 $\pm$ 0.0066)% $Z \rightarrow \tau^+ \tau^-$ (3.3696 $\pm$ 0.0083)% $Z \rightarrow \ell^+ \ell^- \ell^+ \ell^-$ (3.5 $\pm$ 0.4) $\times 10^{-6}$ % $Z \rightarrow \nu \nu$ (20.000 $\pm$ 0.055)% $Z \rightarrow \text{hadrons}$ (69.911 $\pm$ 0.056)%
Cross Sections $\times$ Branching fraction for pp collisions at $\sqrt{s}=13$ TeV	$\sigma(pp \rightarrow W^+ X) \times \mathcal{B}(W^+ \rightarrow \ell^+ \nu) =$ 11370 ± 50 (stat) ± 230 (syst) ± 550 (lumi) pb $\sigma(pp \rightarrow W^- X) \times \mathcal{B}(W^- \rightarrow \ell^- \bar{\nu}) =$ 8580 ± 50 (stat) ± 160 (syst) ± 410 (lumi) pb	$\sigma(pp \rightarrow ZX) \times \mathcal{B}(Z^+ \rightarrow \ell^+ \ell^-) =$ 1910 ± 10 (stat) ± 40 (syst) ± 90 (lumi) pb

in proton-proton collisions. The production of events with W +jets is an important background for the search of SUSY VBF with one muon in the final state. The neutrinos produced in the W decay generate E_T^{miss} , the muon is obtained from the leptonic decay of the W . It is also possible to have fake detector signals identified as muons. On the other hand, it is possible to have jets similar to the VBF jets coming from initial state radiation (ISR) or final state radiation (FSR).

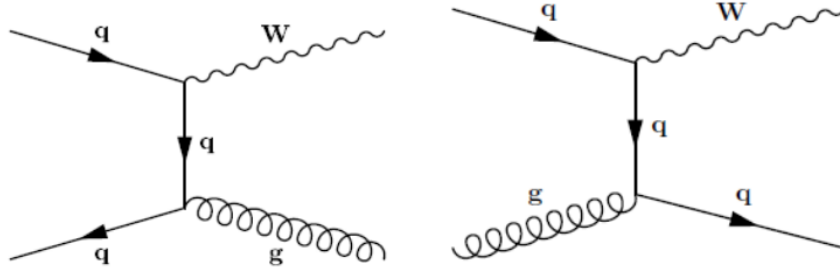


Figure 1.8: Typical diagrams for W production, [62].

1.6.2 Top quark

Table 1.4 displays the world average measurements for mass, decay width, spin, and branching fraction of the top quark. Table 1.4 also shows the cross sections measured by the CMS experiment at $\sqrt{s} = 13$ TeV for $t\bar{t}$ production [63]. The quark top is the only one that experiences a semi-weak decay ($t \rightarrow Wb$) because it is heavier than the W boson. In addition, its lifetime is too short to produce hadronization. A fundamental parameter of the SM is the mass of the top quark (m_{top}) [66]. Providing a precise measurement of the quark top properties is one way to test the internal consistency of the SM, because it allows determination of the couplings for the electroweak and strong interaction. [67, 68]. The most precise measurements of m_{top} have been derived from combinations of the results from the CDF and D0 experiments at the Tevatron, and ATLAS and CMS at the CERN LHC. The current combination from the four experiments gives a top quark mass of 173.34 ± 0.76 GeV [69], a measurement with 0.44% uncertainty. Figure 1.9 shows a summary of the different measurements of the top quark mass compared to the Tevatron measurements and the world average. An additional result that shows the excellent agreement between the experimental results and the theoretical predictions of the SM is the $t\bar{t}$ production cross section as function of the center of mass energy. Figure 1.10 shows the most accurate CMS measurements in the dilepton and lepton + jets channel in comparison with a theory calculation at NNLO + NNLL accuracy.

$t\bar{t}$ production

As shown in the diagrams of Figure 1.11, the top quark decays mainly in a W -boson and a b -quark. $t\bar{t}$ events are the main backgrounds that mimic the final states of the SUSY signal we are searching for. It is possible to obtain a similar signal to the VBF jets due to the misidentification of the two b -quarks or the hadronic decay of the W . Also it is possible to have ISR jets that resemble the VBF jets. E_T^{miss} is also obtained due to the neutrinos produced in the leptonic decays of the W -boson. Finally, a muon can be generated in the decay of the W , or due to fake muons produced after the hadronization.

Table 1.4: Average world measurements for mass, width, spin, and branching ratios of the top quark. Values extracted from [13]. Cross sections values measured by the CMS experiment at $\sqrt{s} = 13$ TeV extracted from [63–65].

Properties	Top quark
Mass (direct measurements)	173.0 ± 0.4 GeV
Decay Width	$1.41^{+0.19}_{-0.15}$ GeV
Spin	$\frac{1}{2}$
Decay Modes (Branching Fractions)	$t \rightarrow Wb \rightarrow e\nu b$ (13.3±0.6)% $t \rightarrow Wb \rightarrow \mu\nu b$ (13.4±0.6)% $t \rightarrow Wb \rightarrow \tau\nu b$ (7.1±0.6)% $t \rightarrow Wb \rightarrow q\bar{q}b$ (66.5±1.4)%
Cross Sections for pp collisions at $\sqrt{s} = 13$ TeV	$\sigma(pp \rightarrow t\bar{t})$ 746 ± 58 (stat) ± 53 (syst) ± 36 (lumi) pb [63] $\sigma(pp \rightarrow tq)$ 238 ± 13 (stat) ± 29 (syst) pb [64] $\sigma(pp \rightarrow tW)$ 63.1 ± 1.8 (stat) ± 6.4 (syst) ± 2.1 (lumi) pb [65]

Single top production

Figure 1.12 shows the Feynman diagrams for single top quark production. Like in the previous case, this process also could mimic our final state due to the leptonic decay of the W. Nevertheless, the cross section is ~ 2.5 times smaller than $t\bar{t}$ production (as shown in Table 1.4).

1.6.3 Diboson production (WW, WZ, ZZ)

Diboson production may have final states similar to the topology of our signal. However, these processes have lower cross sections compared to $t\bar{t}$ or W boson production. Figure 1.13 shows leading order Feynman diagrams for WZ production. Figure 1.14 are the corresponding LO diagrams for WW and ZZ production. Table 1.5 lists measurements of diboson production cross sections at $\sqrt{s} = 13$ TeV.

Table 1.5: Diboson production cross sections measured by the CMS experiment at $\sqrt{s} = 13$ TeV for pp collisions.

Process	Cross Sections for pp collisions at $\sqrt{s} = 13$ TeV
$\sigma(pp \rightarrow WZ)$	$48.09^{+1.00}_{-0.96}$ (stat) $^{+0.44}_{-0.37}$ (theo) $^{+2.39}_{-2.17}$ (syst) ± 1.39 (lumi) pb [73]
$\sigma(pp \rightarrow WW)$	115.3 ± 5.8 (stat) ± 5.7 (syst) ± 6.4 (theo) ± 3.6 (lumi) pb [74]
$\sigma(pp \rightarrow ZZ)$	17.2 ± 0.5 (stat) ± 0.7 (syst) ± 0.4 (theo) ± 0.4 (lumi) pb [75]

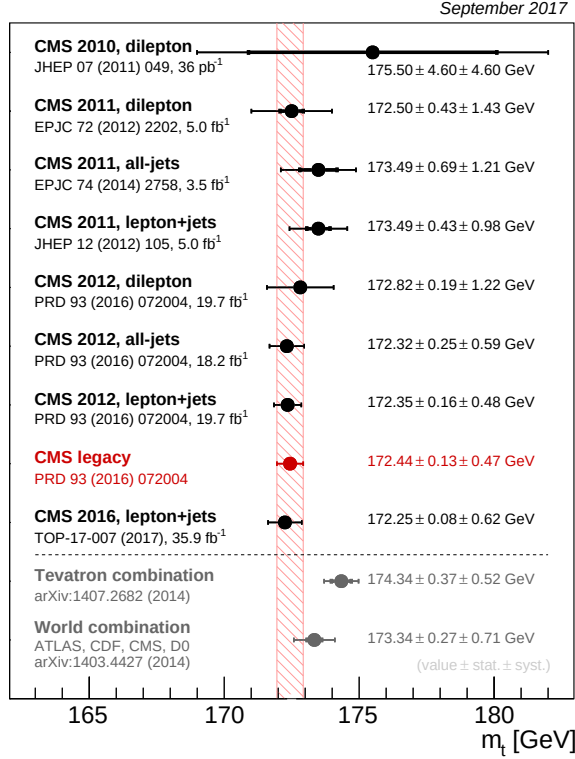


Figure 1.9: Summary of the CMS m_{top} measurements and their combination. The thick error bars show the statistical uncertainty and the thin error bars show the total uncertainty. Also shown are the current Tevatron [70] and world average [69] combinations. Extracted from [71].

1.7 Problems of the Standard Model

No evidence of new physics beyond the SM has been found up to now. Furthermore, the SM has proven to be a successful theory, provided the high precision measurements accumulated and now confirmed at the LHC. Nevertheless, there are important experimental and theoretical reasons that lead us to conclude it is not the final theory to explain nature, but just an effective theory that works pretty well at low energies. Some of the experimental evidences that indicate failures of the SM are:

- **Dark matter and energy:** In 1933 Fritz Zwicky applied the virial theorem to the Coma Cluster and found a discrepancy between the estimated mass (based on the motions of galaxies near its edge) and the expected mass (based on the brightness and number of galaxies) [78]. The conclusion from Zwicky was that to hold the cluster together some unseen matter provided the mass and associated gravitational attraction, providing the first formal inference about the dark matter existence. Then, between the 60's and 70's further strong evidence was provided by Vera Rubin and Kent Ford also using galaxy rotation curves [79,80]. Additional support came from the observations of anisotropies in the cosmic Microwave Background (CMB). The Planck experiment (successor of COBE [81] and WMAP [82] experiments) provided measurements

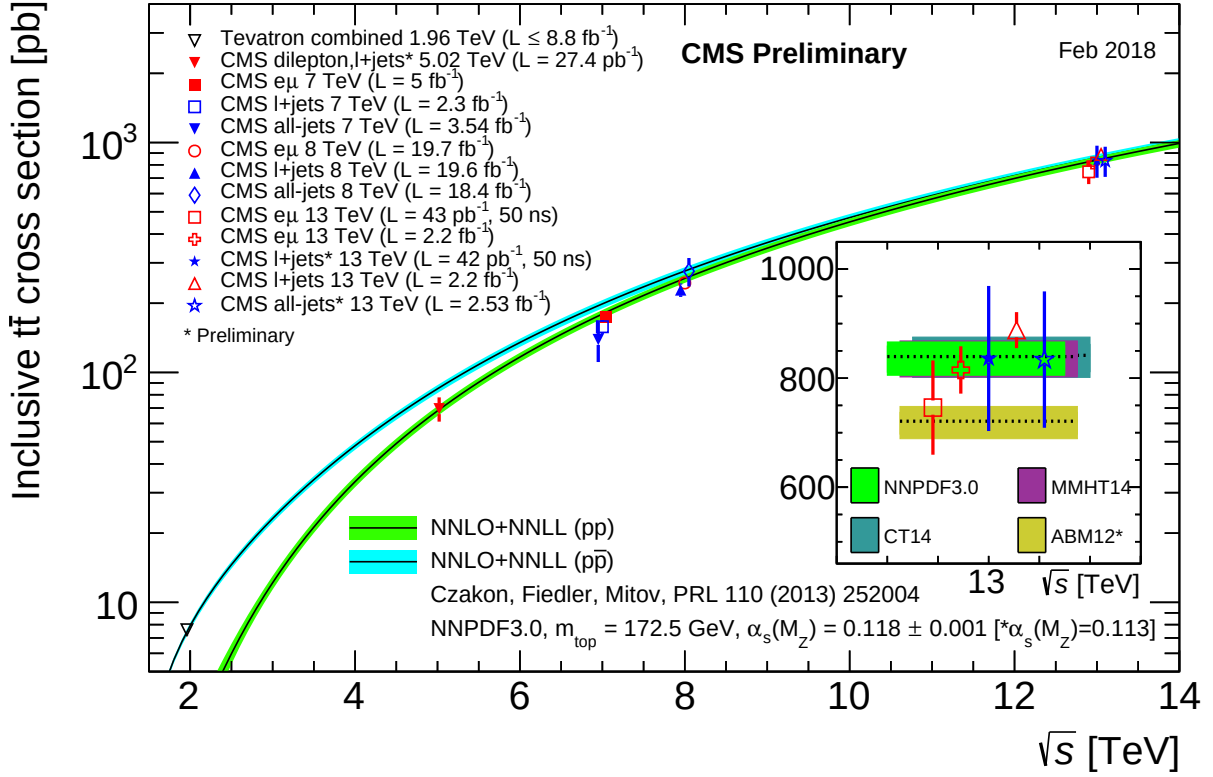


Figure 1.10: $t\bar{t}$ production cross sections measurements for different center of mass energies, performed by the CMS experiment, in the dilepton and lepton+jets channels, in comparison to theory calculation at NNLO+NNLL accuracy. Tevatron measurements are also shown [71].

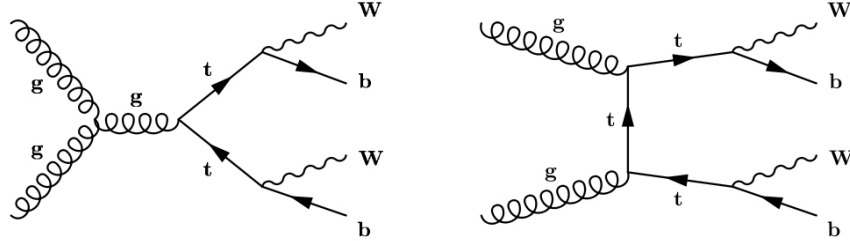


Figure 1.11: Typical diagrams for the $t\bar{t}$ production, extracted from [62].

that are in good agreement with the Lambda cold dark matter model (Λ CDM) [83]. According to this model, the universe contains 68.3% of dark energy, 26.8% of dark matter and only 4.9% ordinary matter and energy. The SM does not provide any dark matter candidate, and also does not provide any explanation for dark energy, failing to explain 95.1% of the total mass-energy content of our universe.

- Gravity: One of the most important recent discoveries is the observation of gravitational

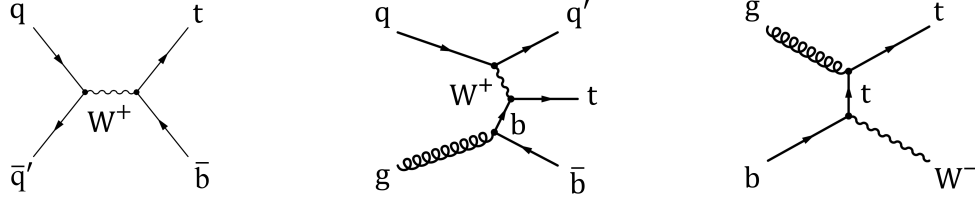


Figure 1.12: Leading-order Feynman diagram for single top quark production. The s channel (left), the dominant next-to-leading-order diagram in the t channel (middle), and the tW production channel (right) [72].

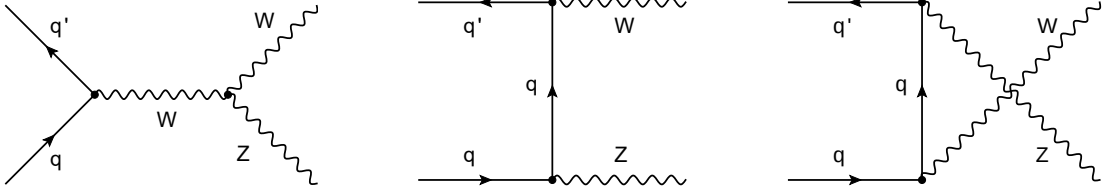


Figure 1.13: Leading-order Feynman diagrams for WZ production in proton-proton collisions. The three diagrams represent contributions from (left) s -channel, (middle) t -channel, and (right) u -channel [76].

waves announced by LIGO and Virgo collaborations on February 2016 [84, 85], a successful prediction from general relativity [86–88]. Nevertheless, the SM does not contain gravitational interactions, and currently there is not a quantum renormalizable field theory for gravity [89].

- **Neutrino masses:** In principle neutrinos are massless particles according to the SM. However, the observation of neutrino oscillations [90, 91] shows that neutrinos have mass. Mass terms for the neutrinos can be included in the SM by means of a Yukawa coupling term, however, the mass terms have to be extraordinarily small. The models that try to explain this apparent hierarchy, go beyond the SM.
- **Matter-antimatter asymmetry:** There is no mechanism in the Standard Model that explains why the universe is made out of mostly matter. The standard model predicts that the universe should be neutral, matter and antimatter should have been created in equal amounts [92].

There are also theoretical drawbacks of the SM:

- **Hierarchy problem:** The SM is unable to prevent scalar particles from acquiring large masses through radiative corrections [93]. This implies that the Higgs mass receives enormous quantum corrections due to virtual effects from every particle coupled to the Higgs field. Due to these corrections the Higgs mass can be written as [94]:

$$m_{h_{SM}}^2 = (m_h^2)_0 + \Delta M_H^2 \quad (1.55)$$

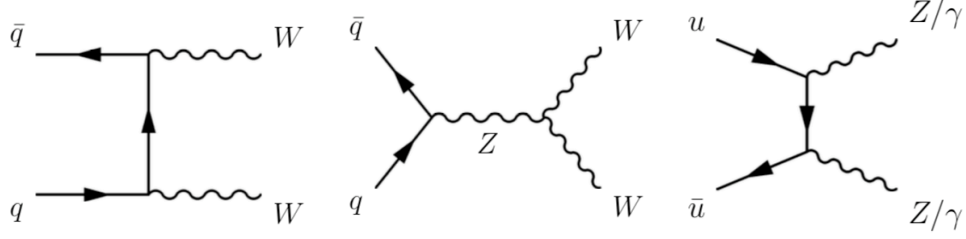


Figure 1.14: Leading-order Feynman diagrams for WW and ZZ production in proton-proton collisions [77].

The term $(m_h^2)_0$ is the bare mass of the Higgs, and ΔM_H^2 is the correction term:

$$\Delta M_H^2 = -\frac{\lambda_f^2}{16\pi^2} \left[2\Lambda^2 + o\left(m_f^2 \ln\left(\frac{\Lambda}{m_f}\right)\right) \right] \quad (1.56)$$

λ_f is the Yukawa coupling constant of the fermions and Λ is the energy scale at which the SM is no longer a valid description of nature. If the SM were to describe nature up to the Planck scale, then the quantum correction ΔM_H^2 would be nearly 30 orders of magnitude larger than the square of the bare Higgs mass (quadratically divergent). This discrepancy between the parameters of the Lagrangian is commonly known as the hierarchy problem [7–9].

- **Grand Unification:** As was discussed in Section 1.5, electromagnetism and the weak interaction are unified in the electroweak interaction of the SM. This unification drives the idea that it is also possible to have the unification of the three interactions of the standard model into a single one at certain high energy threshold [95]. However, it is not possible to obtain any unification from only the evolution of the SM couplings [96]. In order to obtain a unification scenario new physics between the electroweak and the Planck scale is needed.

Chapter 2

Supersymmetry

Supersymmetry (SUSY) is a theoretical symmetry that connects the properties of bosons and fermions [97]. Although its presence has not been verified yet, it is an essential part of many theoretical models, including superstring theory [98]. One of the most exciting predictions in some SUSY models is the existence of a canonical candidate for dark matter, referred to as the Lightest Supersymmetric Particle (LSP). Supersymmetry extends the number of elementary particles so that each SM particle possesses a supersymmetric partner. Superpartners of fermions are bosons, they take names that begin with the letter s, so the selectron is the superpartner of the electron, and the squarks are the superpartners of quarks. Superpartners of bosons are fermions with names ending with -ino, so the photino is the superpartner of the photon, and the gluino is the superpartner of the gluon. SUSY has a high potential to resolve many of the questions of particle physics and astrophysics, as well as solving many theoretical problems of the SM, such as the hierarchy problem, and providing an explanation to Dark Matter [10, 23], this is why it is one of the most attractive models among physicist searching for experimental evidence of physics beyond the standard model at the LHC, with no positive result up to now.

2.1 SUSY Motivations

2.1.1 Unification of Gauge couplings

A strong motivation for SUSY is coming from Grand Unified Theories (GUTs) that state the three gauge interactions from the SM are merged into one at a high energy scale [95]. The first impression is that it is not feasible due to the significant differences presented in the coupling values of the strong, weak and electromagnetic interactions. Nevertheless, the evolution (or running) of these couplings to high energy scales reduces these differences. The couplings depend on energy, an experimental fact that is shown in Figure 2.1 for the strong interaction coupling. According to the renormalization group equations, the strong and weak couplings decrease while the electromagnetic coupling increases with energy [11]. This progression of the couplings makes it probable that at some energy range the couplings could become equal. With that idea, the GUTs propose that this is not accidental, but it is a consequence of the fact that the three interactions possibly have a common origin at some point [11]. However, that unifying interaction splits into three branches because of spontaneous symmetry breaking. Considering that there are precise measurements of the interaction couplings, it is possible to perform a numerical estimation of the unification [96]. Figure 2.2 shows

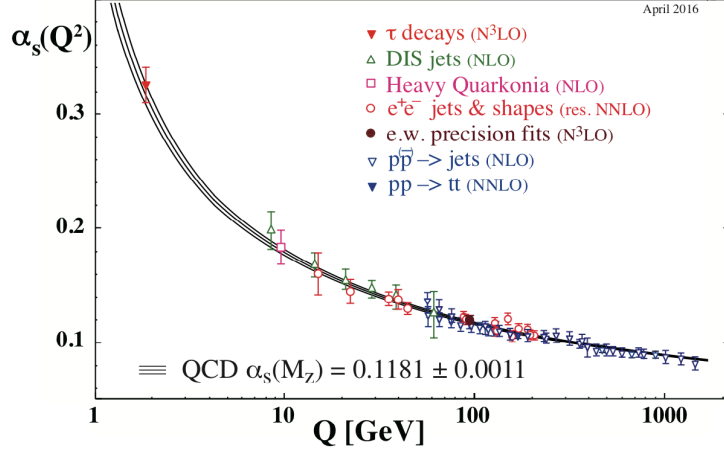


Figure 2.1: Energy dependence of the strong coupling (α_s) extracted from [23].

that it is impossible to achieve any unification by just evolving the SM couplings. This lack of a unifying scheme demands the presence of new physics between the electroweak and the Planck scale to get a unification scenario. Figure 2.2 also shows that SUSY unifies the gauge interactions at a high energy scale. In contrast to the SM, the Minimal Supersymmetric Standard Model (MSSM) [99], the minimum extension of the standard model that includes supersymmetry, has just the right particle content to ensure that it can unify the gauge couplings at a scale of $\sim 1.5 \times 10^{16}$ GeV [21].

2.1.2 Solution to the hierarchy problem

The SM is unable to prevent the Higgs to acquire a quadratic divergent mass through radiative loop corrections. One way of achieving the cancellation of the quadratic divergent terms is supersymmetry. SUSY automatically cancels quadratic corrections in all orders of energy. This cancellation is due to the contributions of the superpartners of ordinary particles. Figure 2.3 shows that the contributions from boson loops cancel those from the fermion ones because of an additional (-1) factor coming from Fermi statistics [11]. Similarly, the contributions from gauge boson loops get canceled out by adding the new gaugino loops. This cancellation is valid in the case of unbroken supersymmetry (partners and superpartners with identical masses) due to the following sum rule relating masses of SM particles and superpartners:

$$\sum_{bosons} m^2 = \sum_{fermions} m^2 \quad (2.1)$$

However, we know that SUSY is a broken symmetry because the different experimental searches for superpartners have excluded already the possibility of having the same mass. In this case, the cancellation is correct up to the SUSY breaking scale, MSUSY, since

$$\sum_{bosons} m^2 - \sum_{fermions} m^2 = M_{SUSY}^2, \quad (2.2)$$

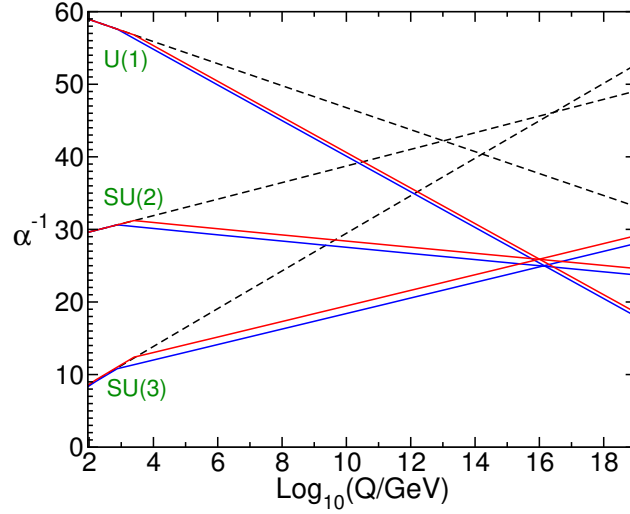


Figure 2.2: Evolution of the inverse gauge couplings $\frac{1}{\alpha}$ for $U(1)_Y$, $SU(2)_L$ and $SU(3)_C$ in the Standard Model (dashed lines) and MSSM scenarios (solid lines) approaching the grand unified theory (GUT) scale. The SUSY particles are assumed to contribute above 750 GeV (blue line) and 2.5 TeV (red line), which causes a change in the slope in the coupling evolution. In addition, the coupling constant for the strong interaction $\alpha_3(m_Z)$ is varied between 0.117 (red line) and 0.120 (blue line). Extracted from [21].

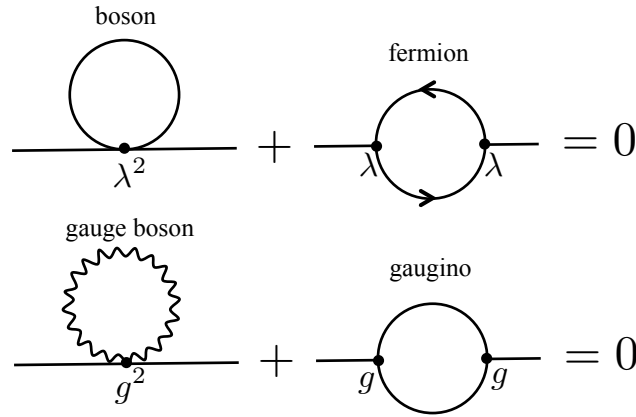


Figure 2.3: Cancellation of quadratic terms (divergencies)

which is expected to be within the 1 TeV scale to make the fine-tuning cancellation of fermion and boson loops natural.

2.1.3 Dark matter in the Universe

According to recent data from the Planck collaboration [83], the matter content of the Universe is the following:

$$\Omega_b = 0.0484 \pm 0.0004 \quad \Omega_{CDM} = 0.258 \pm 0.002$$

Where Ω_b represents the baryon density of the universe and Ω_{CDM} represents the cold dark matter density. The density values show that dark matter exceeds the visible baryonic matter by a factor of 5. One possible explanation for this vast amount of dark matter is hot DM, composed of neutrinos. However, neutrinos are too light (less than about 1 eV), to produce enough DM. On the other hand, they are hot (too fast) to form the dark matter structures necessary to maintain the galaxies and clusters of galaxies together. Another option is cold DM composed of weakly interacting massive particles (WIMPs). However, the standard model does not provide any candidate for cold DM. In counterpart, SUSY provides an excellent candidate called the neutralino, which is the Lightest supersymmetric particle in different SUSY scenarios. Since these hypothetical particles are stable, the neutralinos produced since the Big Bang could have survived [10].

2.2 SUSY Operator Algebra

The SUSY Algebra generator $\hat{Q}$ transforms fermionic states into bosonic states and viceversa as follows:

$$\hat{Q}|\text{fermion}\rangle \propto |\text{boson}\rangle \quad \text{and} \quad \hat{Q}|\text{boson}\rangle \propto |\text{fermion}\rangle \quad (2.3)$$

It changes the spin of the particle by $1/2$. An example of the SUSY generator acting subsequently on a particle of spin 2 gives the following chain of states:

$$\text{spin } 2 \rightarrow \text{spin } 3/2 \rightarrow \text{spin } 1 \rightarrow \text{spin } 1/2 \rightarrow \text{spin } 0.$$

The SUSY generator $\hat{Q}$ satisfies the commutator relations [11]:

$$\{\hat{Q}_\alpha, \hat{Q}_\beta\} = 0 \quad \text{and} \quad \{\hat{Q}_\alpha, \bar{\hat{Q}}_\beta\} = 2\sigma^\mu_{\alpha\beta} P_\mu, \quad (2.4)$$

where P_μ is the momentum four vector and σ^μ are the Pauli matrices. The MSSM has only one generator $\hat{Q}$, so it is a SUSY model with $N = 1$ (the variable N counts the number of SUSY generators). Theories with more than one supersymmetry generator are known as extended supersymmetric theories ($N = 2, 4, 8$) [100]. The MSSM assigns a supersymmetric partner for each particle of the standard model with the same quantum numbers except for spin, which differs by $1/2$.

2.3 Minimal Supersymmetric Standard Model

The MSSM keeps the same gauge symmetries of the Standard Model ($SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$) but doubles the spectrum of new particles. The lack of a mechanism to break the supersymmetry introduces a large number of free parameters in the model (105 parameters) [21, 23, 101]. Theorists initially proposed the MSSM to stabilize the weak scale, solving the hierarchy problem. In the MSSM,

the Higgs boson has a fermionic superpartner: the Higgsino, which would have the same mass of the SM Higgs if there is no symmetry breaking. Since fermion masses are stable to radiative corrections (logarithmically divergent), through this mechanism, the Higgs mass inherits this stability. However, the MSSM would require more than one Higgs field [23]. The MSSM imposes R parity conservation to explain the stability of the proton. The R parity conservation allows the LSP to be stable and weakly interacting. It means that the LSP does not have strong nor electromagnetic interactions, which makes it an excellent DM candidate. The model includes supersymmetry breaking introducing symmetry-breaking operators explicitly on the Lagrangian. The literature has proposed several ways of breaking the symmetry [23], the most common are:

- Gravity mediated [102, 103]
- Gauge mediated [104, 105]
- Anomaly mediated [106, 107]

One of the most popular models for SUSY searches in hadron colliders is the constrained MSSM (CMSSM) [102, 108, 109], also known as minimal supergravity, or MSUGRA. This model reduces a large number of parameters from the MSSM assuming that many of them are equal at large energy scales. The CMSSM considers that the unification takes place at the M_{GUT} , diminishing the number of free parameters to five [23]:

- Defined at the GUT scale:
 - $m_0 \equiv$ Common sfermion mass
 - $m_{1/2} \equiv$ Common gaugino mass
 - $A_0 \equiv$ Common trilinear coupling parameter
- Defined at the electroweak scale:
 - $\tan \beta \equiv$ Ratio of the vacuum expectation values of the Higgs fields for up-type and down-type fermions
 - $\text{Sign}(\mu) \equiv$ Sign of the higgsino mass parameter

However, colored SUSY (squarks and gluinos) have production cross sections that are several orders of magnitude larger than the production cross sections for sleptons or charginos as shown in Figure 2.4. Thus, the searches from ATLAS and CMS have restrained the parameter space of constrained SUSY models that are driven by first and second generation squark production together with gluino production. The exclusion limits for constrained supergravity models for squark and gluino masses are beyond 1 TeV, imposing tight constraints on the allowed parameter space. In recent years, as a consequence of the exclusion limits, less-constrained SUSY models have complemented these constrained models, like the phenomenological-MSSM, or pMSSM [110–112], and also simplified models [113–116]. Simplified models open the possibility of varying masses and other parameters freely, assuming a limited set of production and decay modes for SUSY particles. One specific example is the compressed spectra scenario where the LSP mass is much closer to squark and gluino masses than the predicted by the CMSSM [117].

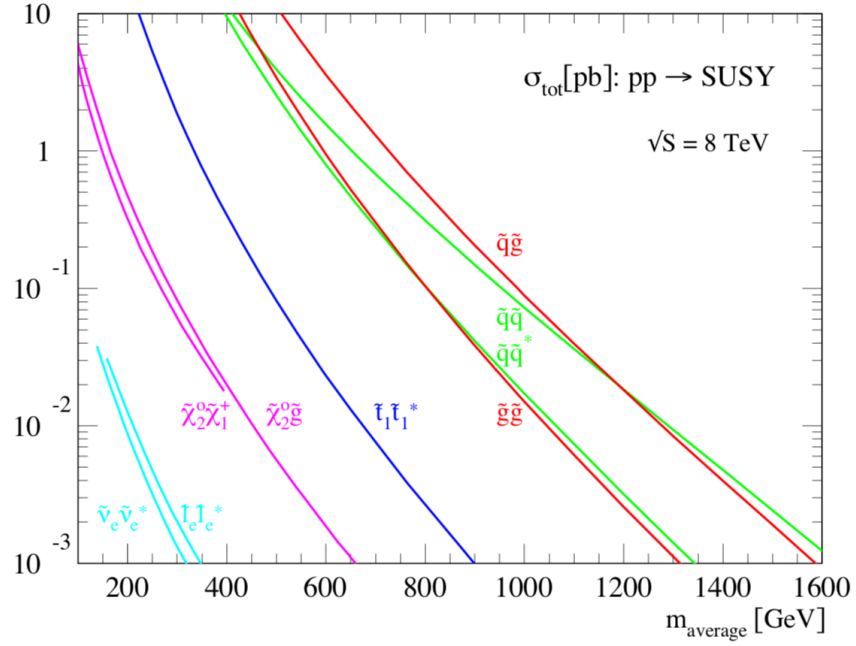


Figure 2.4: Production cross sections for pairs of different sparticles as a function of their mass at the LHC for a center-of-mass energy of 8 TeV. The cross section for the production of squarks assumes mass degeneracy of left and right-handed u, d, s, c and b squarks, with the exception of pair production of stops (explicitly shown). Commonly, colored susy (squarks and gluinos) have production cross sections that are several orders of magnitude larger than the production cross sections for sleptons or charginos [23].

2.4 Lagrangian of the MSSM

The Lagrangian of the MSSM is made of two parts:

$$\mathcal{L} = \mathcal{L}_{SUSY} + \mathcal{L}_{Breaking} \quad (2.5)$$

The first term $\mathcal{L}_{SUSY}$ describes the supersymmetric generalization of the SM. But as the symmetry of the SM particles and their superpartners is broken, then the Lagrangian must contain a term corresponding to the symmetry breaking mechanism [11]. In this case is $\mathcal{L}_{Breaking}$. The supersymmetric term $\mathcal{L}_{SUSY}$ includes kinetic terms that are gauge invariant corresponding to gauge groups $SU(3)$, $SU(2)$, $U(1)$ which depend on three gauge couplings as in the Standard Model, and a supersymmetric extension of the Yukawa term that describes the interactions of the particles with each other:

$$\mathcal{L}_{SUSY} = \mathcal{L}_{Gauge} + \mathcal{L}_{Yukawa} \quad (2.6)$$

the first term is given by

$$\begin{aligned} \mathcal{L}_{Gauge} = & \sum_{SU(3), SU(2), U(1)} \frac{1}{4} \left(\int d^2\theta \, Tr W^\alpha W_\alpha + \int d^2\bar{\theta} \, Tr \bar{W}^{\dot{\alpha}} \bar{W}_{\dot{\alpha}} \right) \\ & + \sum_{Matter} \int d^2\theta d^2\bar{\theta} \, \Phi_i^\dagger e^{g_3 \hat{V}_3 + g_2 \hat{V}_2 + g_1 \hat{V}_1} \Phi_i \end{aligned} \quad (2.7)$$

and the second term is given by

$$\mathcal{L}_{Yukawa} = \int d^2\theta \, (\mathcal{W}_R + \mathcal{W}_{NR}) + h.c. \quad (2.8)$$

where

- θ and $\bar{\theta}$ are Grassmannian anticommuting variables

$$\{\theta_\alpha, \theta_\beta\} = 0, \quad \{\bar{\theta}_{\dot{\alpha}}, \bar{\theta}_{\dot{\beta}}\} = 0, \quad \theta_\alpha^2 = 0, \quad \bar{\theta}_{\dot{\alpha}}^2 = 0, \quad \alpha, \beta, \dot{\alpha}, \dot{\beta} = 1, 2.$$

used to define the superspace that differs from the ordinary Minkowski space:

$$\begin{array}{ccc} Space & \Rightarrow & Superspace \\ x_\mu & & x_\mu, \theta_\alpha, \bar{\theta}_{\dot{\alpha}} \end{array}$$

- W_α is the strength tensor (as analog of $F_{\mu\nu}$ in gauge theories)

$$\begin{aligned} W_\alpha &= -\frac{1}{4} \bar{D}^2 e^V D_\alpha e^{-V}, \\ \bar{W}_{\dot{\alpha}} &= -\frac{1}{4} D^2 e^V \bar{D}_{\dot{\alpha}} e^{-V}, \end{aligned} \quad (2.9)$$

The D s are the supercovariant derivatives given by:

$$\bar{D} = -\frac{\partial}{\partial \bar{\theta}} - i\theta^\mu \partial_\mu \quad (2.10)$$

- g_1, g_2 and g_3 are the usual $U(1), SU(2)$ and $SU(3)$ coupling constants.
- Φ_i are chiral superfields obeying:

$$\bar{D}\Phi_i = 0 \quad (2.11)$$

and V is a real vector superfield $V = V^+$ needed to construct gauge invariant interactions (similar to the vector field A_μ introduced in Section 1.2 to preserve gauge invariance). Φ^+ is an antichiral superfield obeying the equation

$$D\Phi^+ = 0, \quad \text{with } D = \frac{\partial}{\partial\theta} + i\sigma^\mu\bar{\theta}\partial_\mu.$$

The product of chiral superfields Φ^2, Φ^3 , etc is also a chiral superfield. The product of antichiral superfields is also an antichiral superfield, while the product of chiral and antichiral fields $\Phi^+\Phi$ is a general superfield. To obtain a gauge-invariant interaction with matter chiral superfields, we use the gauge transformation:

$$\Phi \rightarrow e^{-ig\Lambda}\Phi, \quad \Phi^+ \rightarrow \Phi^+e^{ig\Lambda^+}, \quad V \rightarrow V + i(\Lambda - \Lambda^+),$$

where Λ is a gauge parameter (chiral superfield). To construct both the SUSY and gauge invariant kinetic term (compared with the covariant derivative in a usual gauge theory) we use:

$$\Phi_i^+\Phi_i \Rightarrow \Phi_i^+e^{gV}\Phi_i \quad (2.12)$$

The superpotential W_R coming from the Yukawa term $\mathcal{L}_{Yukawa}$ looks as follows:

$$W_R = \epsilon_{ij}(h_{ab}^U Q_a^j U_b^c H_2^i + h_{ab}^D Q_a^j D_b^c H_1^i + h_{ab}^L L_a^j E_b^c H_1^i + \mu H_1^i H_2^j), \quad (2.13)$$

where

- the superfields are listed in Table 2.1,
- $i, j = 1, 2$ are the indices of $SU(2)$
- $a, b = 1, 2, 3$ are the generation indices
- $h^{U,D,L}$ are the Yukawa couplings
- ϵ_{ij} is the antisymmetric tensor.

The charge conjugation in the fields of Equation 2.13 is denoted by the letter c . This part of the Lagrangian is similar to the Standard Model, but in comparison the superpotential contains only superfields instead of ordinary SM fields. Additionally, the last term describes the Higgs mixture, which is absent in the SM, as this one has only one Higgs field [10]. The second part of Equation 2.8 is

$$\begin{aligned} W_{NR} &= \epsilon_{ij}(\lambda_{abd}^L L_a^i L_b^j E_d^c + \lambda_{abd}^{L'} L_a^i Q_b^j D_d^c + \mu'_a L_a^i H_2^j) \\ &+ \lambda_{abd}^B U_a^c D_b^c D_d^c. \end{aligned} \quad (2.14)$$

in the SM Lagrangian these terms are not present. Due to the Lorentz invariance, the superfields in Equation (2.14) can not be replaced by ordinary SM fields like in Equation (2.13). The first line in Equation (2.14) violates lepton number and the second line violates baryon number. So these terms have to be suppressed or excluded in some way because lepton and baryon number violation is not observed in nature.

2.4.1 R-Parity

The introduction of a special symmetry called R-parity avoids the lepton and/or baryon number violation [12, 118]. R-parity is a multiplicative quantum number defined as:

$$R = (-1)^{3(B-L)+2S}, \quad (2.15)$$

where B is the baryon number, L is the lepton number and S is the spin of the particle [11]. Therefore R is equal to 1 for all leptons and quarks; and equals to -1 for the SM superpartners due to the factor $(-1)^{2S}$. The W_{NR} terms in Equation 2.14 are banned due to the R-parity, so normally it is assumed that those terms are not present in the MSSM, in other words, that means the preservation of R-parity. Nevertheless, it would be possible that those terms are present, but this will imply that there are very severe experimental limits on the couplings [119–122]:

$$\lambda_{abc}^L, \lambda_{abc}^{L'} < 10^{-4}, \quad \lambda_{abc}^B < 10^{-9}.$$

In conclusion, the superpotential can only contain terms that violate the baryon number conservation or the lepton number conservation. As no violation of these numbers has been observed, these terms should be removed from the superpotential. The conservation of R-parity has the following consequences:

- The LSP is stable, because its decay into ordinary matter would change the R-parity. Additionally, if the LSP particle has no charge, it is a perfect candidate for dark matter. In addition, for the production of supersymmetric particles in hadron colliders, once one LSP particle is produced, it could not decay into known SM particles.
- The decay of SUSY particle products should contain an odd number of LSPs. This is a feature of the production of SUSY matter in the LHC, which could be used to distinguish particles coming from SUSY events from particles coming from SM events.
- The superpartners of the SM must be produced in pairs. Consequently if the R-parity is conserved, all the SUSY events should always produce at least two LSPs.

2.5 Particle spectrum of the MSSM

Table 2.1 shows the particle spectrum of the MSSM. The squarks and sleptons are spin zero particles, then the labels L and R are used to indicate that they are superpartners of left or right-handed particles. Table 2.1 also shows the presence of an additional Higgs boson, for this reason, there are two doublets of scalar fields [10]. Another essential characteristic of the SUSY particle spectrum is mixing, the superpartners of gauge bosons (gauginos) and Higgs (higgsinos) are mixed. Then the observable values, neutralinos, and charginos are formed by combinations of gauginos and higgsinos [23].

2.5.1 Chargino mixing

The charginos are mixed states of higgsinos and gauginos. The matrix of mixing that shows how the composition of the neutralinos is in terms of the higgsinos and gauginos is given by:

Superfield	Bosons		Fermions		$SU_c(3)$	$SU_L(2)$	$U_Y(1)$
Gauge							
$\mathbf{G}^a$	gluon	g^a	gluino	$\tilde{g}^a$	8	0	0
$\mathbf{V}^k$	Weak	W^k ($W^\pm, Z$)	wino, zino	$\tilde{w}^k$ ($\tilde{w}^\pm, \tilde{z}$)	1	3	0
$\mathbf{V}'$	Hypercharge	B (γ)	bino	$\tilde{b}(\tilde{\gamma})$	1	1	0
Matter							
$\mathbf{L}_i$	sleptons	$\left\{ \begin{array}{l} \tilde{L}_i = (\tilde{\nu}, \tilde{e})_L \\ \tilde{E}_i = \tilde{e}_R \end{array} \right.$	leptons	$\left\{ \begin{array}{l} L_i = (\nu, e)_L \\ E_i = e_R \end{array} \right.$	1	2	-1
$\mathbf{E}_i$					1	1	2
$\mathbf{Q}_i$	squarks	$\left\{ \begin{array}{l} \tilde{Q}_i = (\tilde{u}, \tilde{d})_L \\ \tilde{U}_i = \tilde{u}_R \\ \tilde{D}_i = \tilde{d}_R \end{array} \right.$	quarks	$\left\{ \begin{array}{l} Q_i = (u, d)_L \\ U_i = u_R^c \\ D_i = d_R^c \end{array} \right.$	3	2	1/3
$\mathbf{U}_i$					3^*	1	-4/3
$\mathbf{D}_i$					3^*	1	2/3
Higgs							
$\mathbf{H}_1$	Higgses	$\left\{ \begin{array}{l} H_1 \\ H_2 \end{array} \right.$	higgsinos	$\left\{ \begin{array}{l} \tilde{H}_1 \\ \tilde{H}_2 \end{array} \right.$	1	2	-1
$\mathbf{H}_2$					1	2	1

Table 2.1: Particle spectrum of the MSSM [123]. The particles and their superpartners are arranged in super-multiplets, where $a = 1 \dots 8$ is the index of $SU_c(3)$ and $k = 1 \dots 3$ is the index of $SU_L(2)$. The three generations are labeled with the index $i = 1 \dots 3$. Hereafter, the tilde denotes a superpartner of an ordinary particle.

$$\mathcal{M}_C = \begin{pmatrix} M_2 & \sqrt{2}m_W \sin \beta \\ \sqrt{2}m_W \cos \beta & \mu \end{pmatrix}, \quad (2.16)$$

In this matrix, the $SU(2)$ gaugino mass parameter is represented by M_2 (which can be real and positive). On the other hand, the higgsino mass parameter (μ) can have complex values [124]. Additionally, to diagonalize the $\mathcal{M}_C$ matrix two unitary matrices U and V are needed because the $\mathcal{M}_C$ matrix is not symmetric

$$U^* \mathcal{M}_C V^\dagger = \begin{pmatrix} m_{\tilde{\chi}_1^\pm} & 0 \\ 0 & m_{\tilde{\chi}_2^\pm} \end{pmatrix} \quad \text{with} \quad m_{\tilde{\chi}_1^\pm} < m_{\tilde{\chi}_2^\pm}. \quad (2.17)$$

In this way, we get the two mass values $m_{\tilde{\chi}_1^\pm}, m_{\tilde{\chi}_2^\pm}$:

$$m_{\tilde{\chi}_1^\pm, \tilde{\chi}_2^\pm}^2 = \frac{1}{2} \left[M_2^2 + \mu^2 + 2m_W^2 \mp \sqrt{(M_2^2 - \mu^2)^2 + 4m_W^4 \cos^2 2\beta + 4m_W^2 (M_2^2 + \mu^2 + 2M_2\mu \sin 2\beta)} \right]. \quad (2.18)$$

2.5.2 Neutralino mixing

The electroweak gauginos and higgsinos masses are also mixed. These mass matrices can be diagonalized using unitary matrices, where the diagonal elements are the physical masses of the particles,

the neutralinos [124]. The mixing matrix in terms of the neutral higgsinos ($\tilde{H}_1^0$ and $\tilde{H}_2^0$) and the neutral gauginos ($\tilde{B}$ and $\tilde{W}^3$), is given by:

$$\mathcal{M}_N = \begin{pmatrix} M_1 & 0 & -m_Z c_\beta s_W & m_Z s_\beta s_W \\ 0 & M_2 & m_Z c_\beta c_W & -m_Z s_\beta c_W \\ -m_Z c_\beta s_W & m_Z c_\beta c_W & 0 & -\mu \\ m_Z s_\beta s_W & -m_Z s_\beta c_W & -\mu & 0 \end{pmatrix} \quad (2.19)$$

where M_1 and M_2 are the $U(1)$ and $SU(2)$ gaugino masses, μ is the higgsino mass parameter and $\tan \beta = v_2/v_1$ ($c_\beta = \cos \beta$, $s_W = \sin \theta_W$). Unlike the previous case, only one unitary matrix is needed to diagonalize the $\mathcal{M}_N$ matrix, since the matrix is symmetric. The eigenvalues of this matrix give the physical masses of the neutralinos:

$$\text{diag}(m_{\tilde{\chi}_1^0}, m_{\tilde{\chi}_2^0}, m_{\tilde{\chi}_3^0}, m_{\tilde{\chi}_4^0}) = N^* \mathcal{M}_N N^\dagger, \quad (m_{\tilde{\chi}_1^0} < m_{\tilde{\chi}_2^0} < m_{\tilde{\chi}_3^0} < m_{\tilde{\chi}_4^0}). \quad (2.20)$$

Explicit values for the neutralinos masses can be found in Reference [125].

2.6 SUSY Searches

The largest constraints, for several models proposing SUSY in the TeV scale, come from collider experiments. SUSY searches have been conducted since the 80's, beginning with the UA1 [126] and UA2 [127, 128] experiments at the CERN Super Proton Synchrotron. The Large Electron Positron Collider (LEP) [129] was an electron-positron circular collider with $\sqrt{s}$ ranging between 90 and 209 GeV that was operating from 1989 until 2000. In 2001 it was dismantled and the tunnel was re-utilized for the Large Hadron Collider (LHC, See Section 3.1). From the results, coming from the LEP detectors (ALEPH, DELPHI, OPAL, and L3), exclusion limits were obtained for the mass of charginos, sleptons, and third-generation squarks of around 100 GeV [130]. The LEP limits were updated with the results taken with the Tevatron, which produced proton-antiproton collisions at $\sqrt{s}$ of 1.96 TeV until 2011. If SUSY production would have occurred at the Tevatron, it would have happened strongly in a $q\bar{q}$ interaction since the Tevatron was a $p\bar{p}$ machine. As a result, the CDF and D0 Tevatron experiments set higher limits than the results produced by the LEP experiments in the squark and gaugino productions. Under specific SUSY model conditions, the third-generation squark masses were ruled out up to ~ 240 GeV and the chargino masses up to 170 GeV [131]. Between 1992 and 2007, there was a search of single SUSY particles violating R-parity in the H1 and ZEUS experiments from the Hadron-Electron Ring Accelerator (HERA) [132], an electron-proton collider with $\sqrt{s}$ of 318 GeV. The LHC took data for first time between 2010 and 2012 at $\sqrt{s} = 7$ TeV and 8 TeV, respectively. During that period 20 fb $^{-1}$ of data were accumulated. SUSY parameter regions where potential discoveries were expected to occur were rapidly excluded by ATLAS and CMS experiments. Exclusion limits of ~ 700 GeV were achieved for some scenarios (as shown in Figure 2.5) only with the data taken in 2010 (with just 35 pb $^{-1}$).

2.7 Colored SUSY searches at the LHC

All the initial SUSY searches at the LHC have been focused on the production of heavy colored SUSY particles through strong interactions. The cross sections for these processes are comparable to the

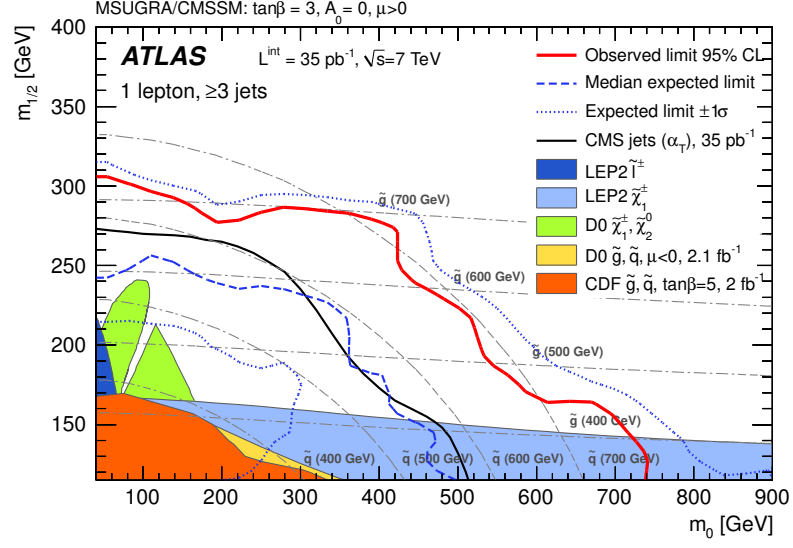


Figure 2.5: Observed and expected exclusion limits at 95% C.L. as well as the $\pm 1\sigma$ variation on the expected limit, for the combined electron and muon channels from the ATLAS collaboration at $\sqrt{s} = 7$ TeV. The published limits from CMS [133], CDF [134], D0 [135, 136] are also shown with the results from the LEP [137] experiments. Extracted from [138].

ones found in the SM [10]. The decays of squarks and gluinos generally result in a large multiplicity of particles in the final state. This means that long decay chains are expected with the production of several jets and/or leptons making this scenario difficult to study. The main production channels in proton-proton collisions at the energies of 8 or 13 TeV for colored SUSY particles are: the fusion of gluons, resulting in squark pairs or gluino pairs (the case of gluon fusion with creation of gluino is presented in Table 2.2). The strategy of previous searches at the LHC for squarks and gluinos is based on the assumption that the masses of these superpartners are in the region of 1 TeV so that they might be created at the mass shell with a cross section high enough to distinguish them from the background of the SM particles [10]. The main decay modes of superpartners produced by colored SUSY particles (together with their signatures at detector level) are presented in Table 2.3. The CMS and ATLAS experiments have constrained a large number of SUSY scenarios. In particular, the present colored searches conducted by CMS and ATLAS have produced exclusion limits for the masses of squarks and gluinos greater than 1.5 TeV. One important case is shown in Figure 2.6 where 95% CL limits are provided for gluino pair production to the decay chains $\tilde{g} \rightarrow t\bar{t}\tilde{\chi}_1^0$ from the CMS collaboration at $\sqrt{s} = 13$ TeV. Depending on the scenario the gluino masses are excluded until 1.5 or 1.9 TeV. Figure 2.7 shows a summary of exclusion limits obtained for different searches of gluinos, squarks and electroweakinos with 36 fb^{-1} at $\sqrt{s} = 13$ TeV by the CMS experiment [139], similar results have been obtained by ATLAS experiment as shown in Figure 2.8. In the gluino case the exclusion has reached ~ 1.9 TeV and squarks ($\tilde{q}$) have been ruled out up to ~ 1.4 TeV for the first and second generations in certain scenarios.

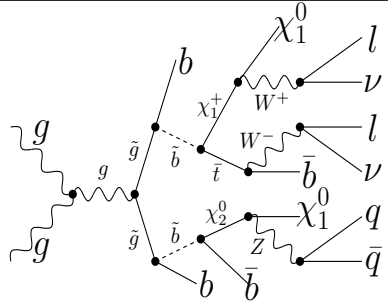
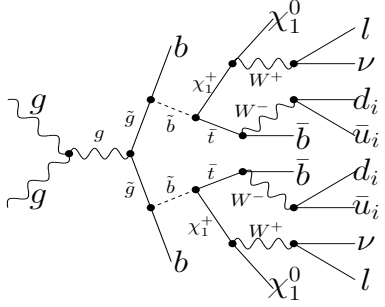
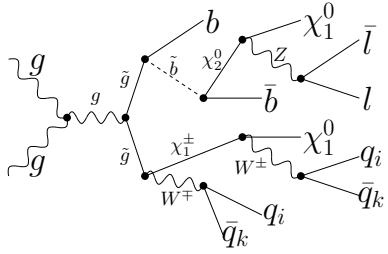
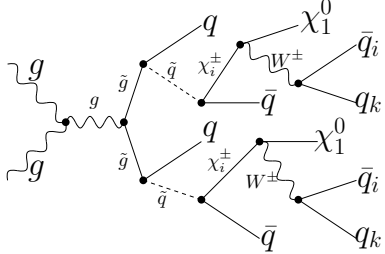
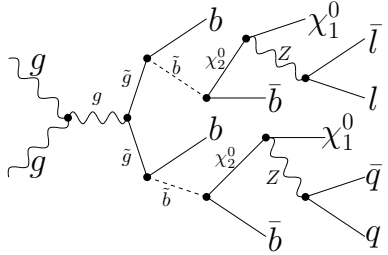
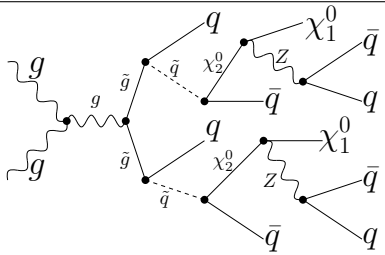
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	2ℓ $6j$ $\cancel{E}_T$		$8j$ $\cancel{E}_T$
	2ℓ $6j$ $\cancel{E}_T$		$8j$ $\cancel{E}_T$

Table 2.2: Creation of the pair of gluino with further cascade decay (extracted from [10])

Table 2.3: Creation of superpartners in several channels with their main decay modes and detector signatures (extracted from [10])

Creation	The main decay modes	Signature
$\tilde{g}\tilde{g}, \tilde{q}\tilde{q}, \tilde{g}\tilde{q}$	$\left. \begin{array}{l} \tilde{g} \rightarrow q\bar{q}\tilde{\chi}_1^0 \\ q\bar{q}'\tilde{\chi}_1^\pm \\ g\tilde{\chi}_1^0 \end{array} \right\} m_{\tilde{q}} > m_{\tilde{g}}$ $\left. \begin{array}{l} \tilde{q} \rightarrow q\tilde{\chi}_i^0 \\ \tilde{q} \rightarrow q'\tilde{\chi}_i^\pm \end{array} \right\} m_{\tilde{g}} > m_{\tilde{q}}$	$\cancel{E}_T + \text{multijets (+leptons)}$
$\tilde{\chi}_1^\pm \tilde{\chi}_2^0$	$\tilde{\chi}_1^\pm \rightarrow \tilde{\chi}_1^0 \ell^\pm \nu, \tilde{\chi}_2^0 \rightarrow \tilde{\chi}_1^0 \ell \ell$ $\tilde{\chi}_1^\pm \rightarrow \tilde{\chi}_1^0 q\bar{q}', \tilde{\chi}_2^0 \rightarrow \tilde{\chi}_1^0 \ell \ell,$	trilepton + $\cancel{E}_T$ dileptons + jet + $\cancel{E}_T$
$\tilde{\chi}_1^+ \tilde{\chi}_1^-$	$\tilde{\chi}_1^+ \rightarrow \ell \tilde{\chi}_1^0 \ell^\pm \nu$	dilepton + $\cancel{E}_T$
$\tilde{\chi}_i^0 \tilde{\chi}_i^0$	$\tilde{\chi}_i^0 \rightarrow \tilde{\chi}_1^0 X, \tilde{\chi}_i^0 \rightarrow \tilde{\chi}_1^0 X'$	dilepton+jet + $\cancel{E}_T$
$\tilde{t}_1 \tilde{t}_1$	$\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0$ $\tilde{t}_1 \rightarrow b\tilde{\chi}_1^\pm, \tilde{\chi}_1^\pm \rightarrow \tilde{\chi}_1^0 q\bar{q}'$ $\tilde{t}_1 \rightarrow b\tilde{\chi}_1^\pm, \tilde{\chi}_1^\pm \rightarrow \tilde{\chi}_1^0 \ell^\pm \nu,$	2 noncollinear jets + $\cancel{E}_T$ single lepton + $\cancel{E}_T + b's$ dilepton + $\cancel{E}_T + b's$
$\tilde{l}\tilde{l}, \tilde{l}\tilde{\nu}, \tilde{\nu}\tilde{\nu}$	$\tilde{\ell}^\pm \rightarrow \ell^\pm \tilde{\chi}_i^0, \tilde{\ell}^\pm \rightarrow \nu_\ell \tilde{\chi}_i^\pm$ $\tilde{\nu} \rightarrow \nu \tilde{\chi}_1^0$	dilepton + $\cancel{E}_T$ single lepton + $\cancel{E}_T$

2.8 Electroweak Searches

In electroweak channels, the direct production of electroweakinos ($\tilde{\chi}_2^0, \tilde{\chi}_1^0, \tilde{\chi}_1^\pm$), expected to be lighter than colored SUSY particles, is allowed. The subsequent direct production of electroweakinos is followed by simple final states with a few leptons and neutralino(s). As a consequence, the decay chains of these channels is simpler than those on the colored sector which makes the searches easier from the experimental point of view. In the case of electroweak SUSY searches, the particle mediators of the process do not have color charge. The same case happens with the SUSY particles produced. All the SUSY particles that satisfy this requirement come from the electroweak sector of the MSSM: vectorial bosons, charginos and neutralinos. In general, the charginos and neutralinos are lighter particles compared with the squarks and gluinos. The exclusion limits on electroweak SUSY searches produced by the ATLAS and CMS collaborations for the Run 1 of the LHC are shown in Figure 2.9. The limits are presented for different $\tilde{\chi}_1^0$ masses as a function of the $\tilde{\chi}_2^0/\tilde{\chi}_1^\pm$ mass. The limits assume scenarios where the mass of the $\tilde{\chi}_2^0$ and $\tilde{\chi}_1^\pm$ are equal, which is allowed by theory since these gauginos belong to the same gauge group multiplet. In comparison with colored searches (described in the previous section), the exclusion limits on the charginos and neutralinos are less constrained, as may be expected in a hadron collider where these particles suffer from smaller electroweak production cross sections. Figure 2.10 show the upper mass limits at 95% C.L. on chargino-neutralino pair production for different decay chains (via sleptons and W/Z/H) from the CMS collaboration at $\sqrt{s} = 13$ TeV. For compressed mass spectra scenarios, where the mass of the second neutralino

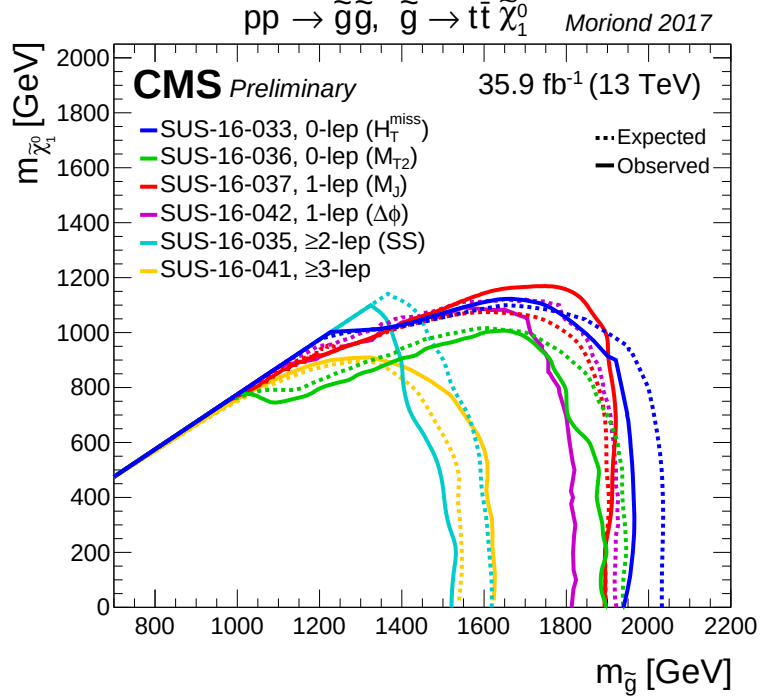


Figure 2.6: Upper mass limits at 95% C.L. on gluino pair production for the decay chains $\tilde{g} \rightarrow t\bar{t}\tilde{\chi}_1^0$ from the CMS collaboration at $\sqrt{s} = 13$ TeV. [14].

$\tilde{\chi}_2^0$, or the mass of the chargino $\tilde{\chi}_1^\pm$, is very close to the sum of the masses of its decay products, the exclusion limits are low. This is because the decay products can only have very low transverse momentum (soft p_T) and are difficult to identify from their corresponding SM backgrounds. So, there is still unexplored territory to probe, in particular, regions with compressed mass spectra scenarios where the sensitivity is reduced.

2.9 Electroweak VBF-SUSY Searches

Given that null results have been obtained in the colored SUSY searches, electroweak SUSY production becomes very relevant. One powerful tool is Vector boson fusion (VBF) tagging, in particular for compressed spectra scenarios, where the mass of the LSP is slightly different than the masses of other charginos and neutralinos [117]. The mechanism of electroweak production includes the emission of vector bosons (W^+ , W^- and Z^0) by the quarks that participate in the collision. These bosons fuse to produce particles from the SM or SUSY spectrum. Figure 2.11 shows the production of two neutralinos with two forward jets through Vector Boson Fusion. This kind of processes suffer from smaller cross sections, but benefit from lower contamination from SM backgrounds due to its singular topological signature of two forward jets with large invariant mass. Besides the large background reduction, the inclusion of the two forward high p_T jets on the event topology allow the analyses to use triggers designed to select events with this type of signature. There are several phenomenological and theoretical papers [142, 143] claiming that compressed spectra scenarios where the $\tilde{\tau}$ and the

Selected CMS SUSY Results* - SMS Interpretation

ICHEP '16 - Moriond '17

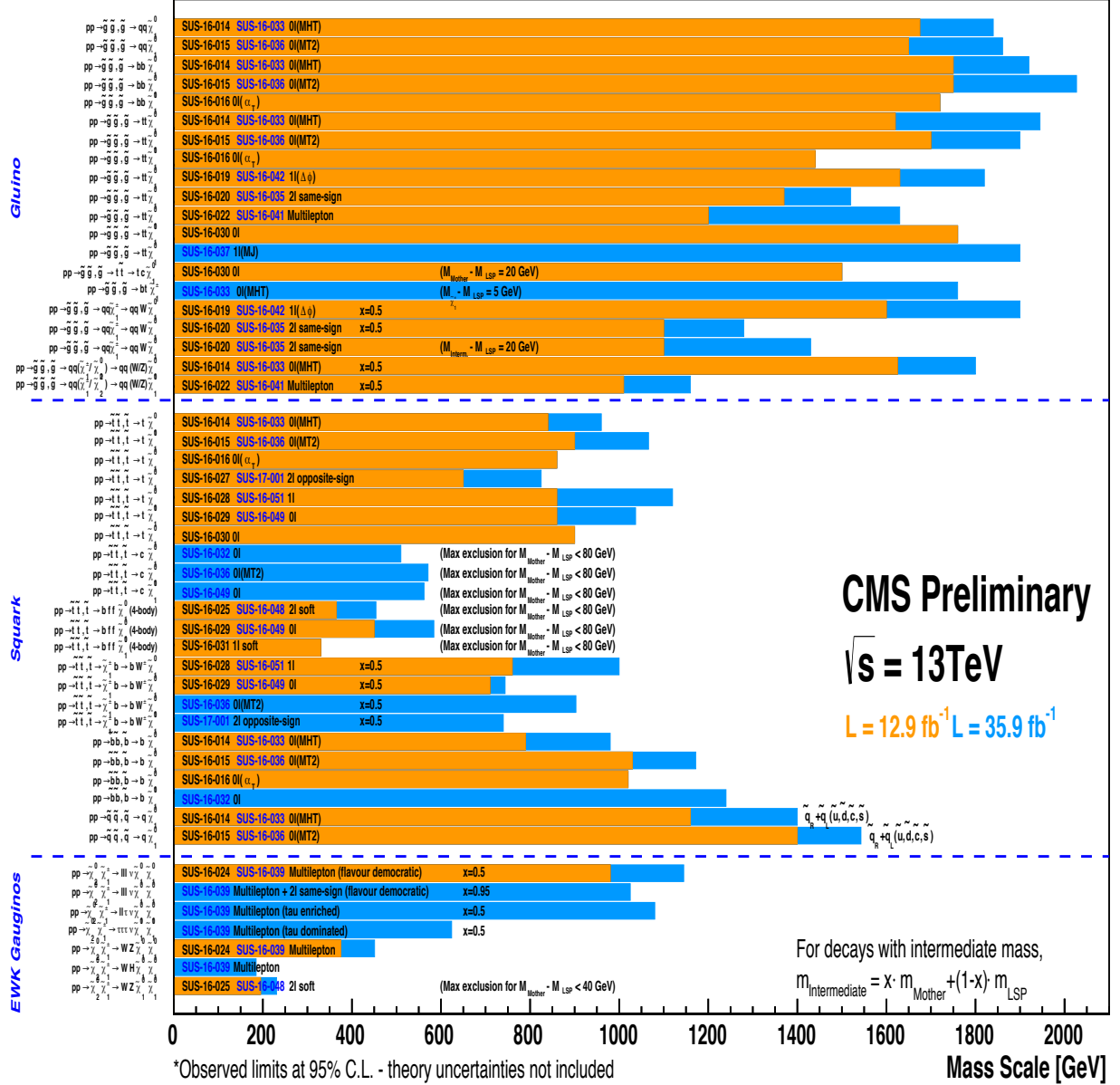


Figure 2.7: CMS summary of SUSY results at $\sqrt{s}$ of 13 TeV [14].

$\tilde{\chi}_1^\pm$ are light and nearly mass degenerate yield the correct dark matter relic density estimated by cosmology. Therefore, VBF SUSY searches become a powerful experimental technique to explore a rich spectrum of SUSY physics models, where classic searches have no experimental reach. Several phenomenological studies [143–145] have proposed the use of dijet events with VBF topology for SUSY searches. The first search, ever, for SUSY using the VBF dijet topology to look for charginos

ATLAS SUSY Searches* - 95% CL Lower Limits

December 2017

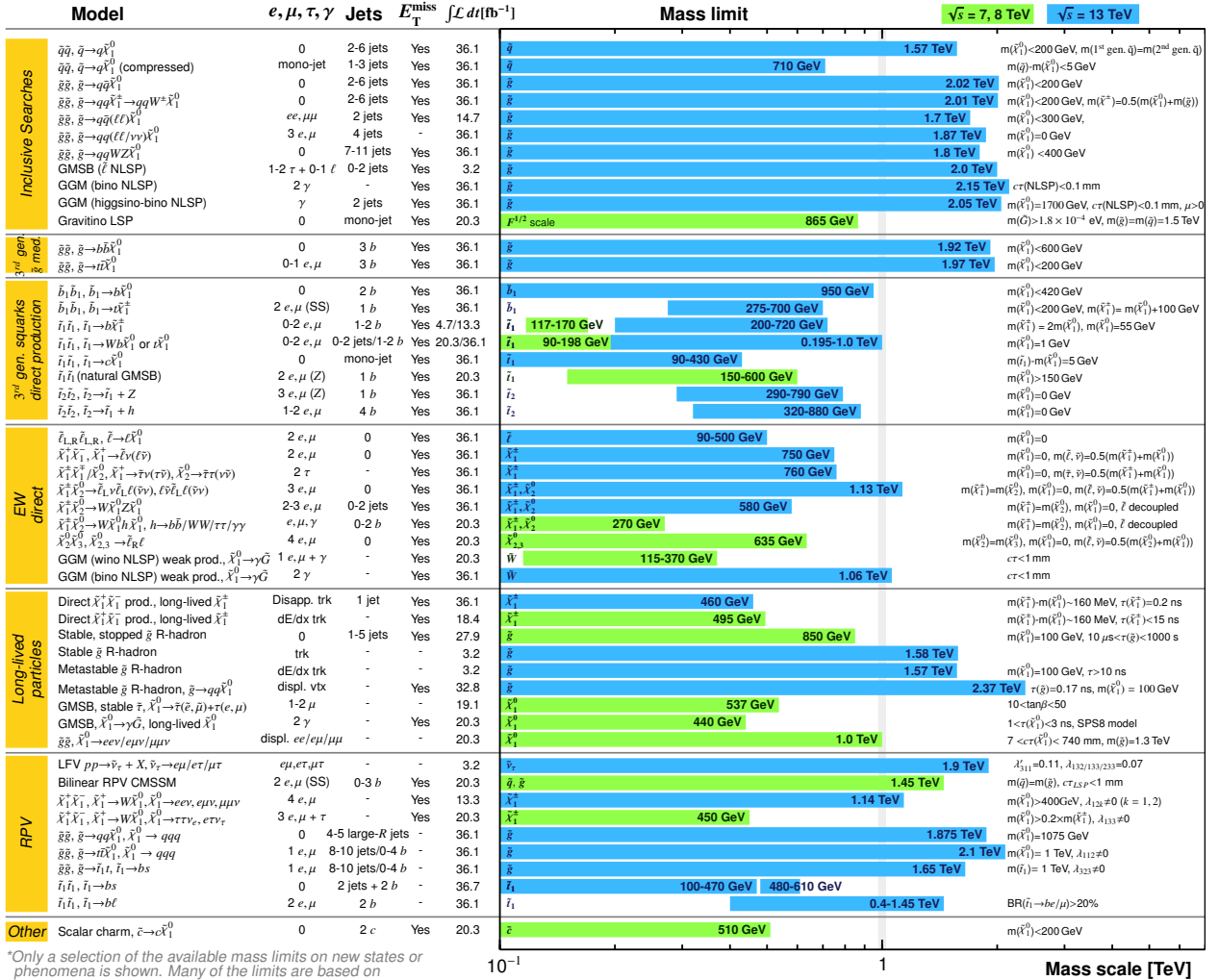


Figure 2.8: Exclusion limits on sparticle masses from SUSY searches by ATLAS collaboration. The 7-8 TeV results are represented by green bars, while the 13 TeV results are summarized by the blue bars. [15].

and neutralinos, using proton-proton collision data at $\sqrt{s} = 8$ TeV (with an integrated luminosity of 19.7 fb^{-1}) is reported in reference [20]. In this search, SUSY models were considered with light third-generation sleptons (staus) leading to leptonic decay modes of the charginos/neutralinos (e.g. $\tilde{\chi}_2^0 \rightarrow \tau \tilde{\tau} \rightarrow \tau \tau \tilde{\chi}_1^0$ as it is shown in the Feynman diagram on the left side of Figure 2.12). In the presence of light sleptons, $\tilde{l} = \tilde{e}, \tilde{\mu}, \tilde{\tau}$, it is likely that $\tilde{\chi}_1^\pm$ as well as $\tilde{\chi}_2^0$ would decay to a $\tilde{l} \nu_l$ and/or a $\tilde{l}^\pm l^\mp$ pair, respectively. Thus, charginos and neutralinos could be searched directly up to a few hundred GeV in final states with leptons and two additional jets according to the VBF topology. The 8 TeV search focused on eight independent channels: $e^\pm \mu^\pm$, $e^\pm \mu^\mp$, $\mu^\pm \mu^\pm$, $\mu^\pm \mu^\mp$, $\mu^\pm \tau_h^\pm$, $\mu^\pm \tau_h^\mp$, $\tau_h^\pm \tau_h^\pm$, and $\tau_h^\pm \tau_h^\mp$. Figure 2.13 (left) shows the expected MC event yield from SM backgrounds in

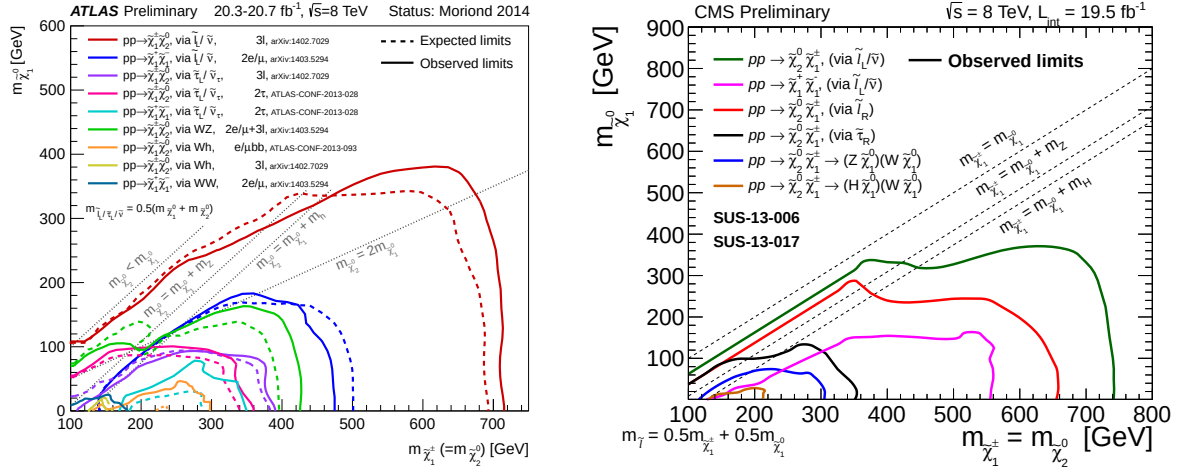


Figure 2.9: (Left) Summary of ATLAS searches for electroweak production of charginos and neutralinos during the Run 1 of the LHC [140]. (Right) Summary of CMS observed limits for EWKino models during the Run 1 of the LHC [141].

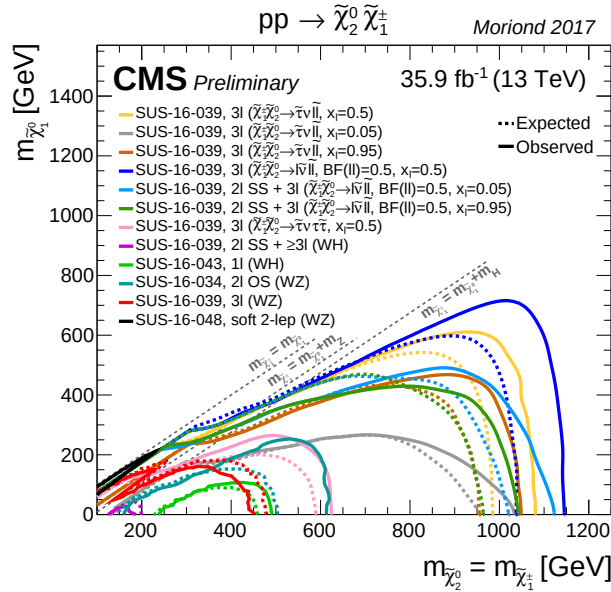


Figure 2.10: Upper mass limits at 95% C.L. on chargino-neutralino pair production for different decay chains (via sleptons and W/Z/H) from the CMS collaboration at $\sqrt{s} = 13$ TeV [14].

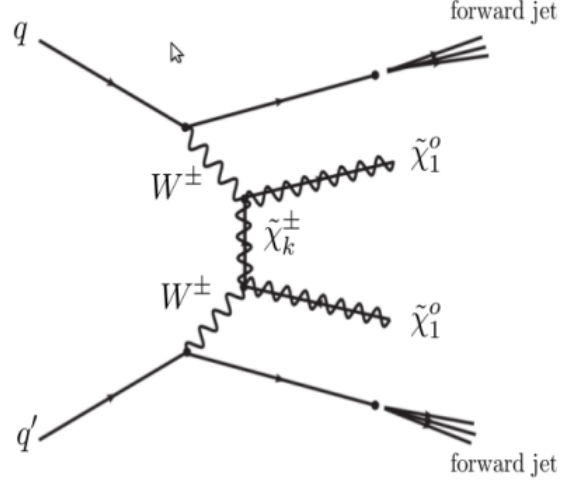


Figure 2.11: Pair production of neutralinos with two forward jets through VBF.

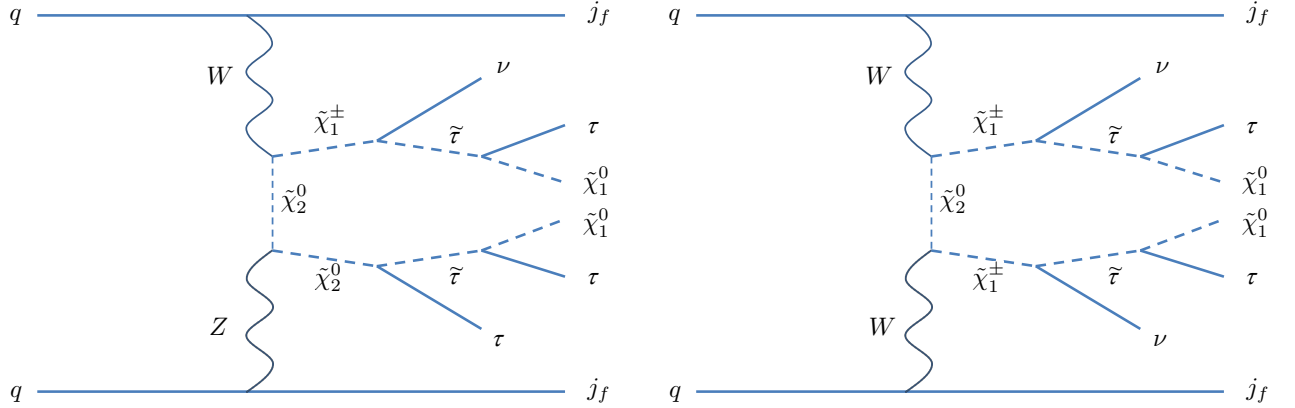


Figure 2.12: Diagrams of (left) chargino-neutralino and (right) chargino-chargino pair production through vector-boson fusion followed by their decays to leptons and the LSP $\tilde{\chi}_1^0$ [20].

the dijet invariant mass distribution at 8 TeV compared to pp collision data. Also, this figure shows the expected SUSY VBF signal (magnified by a factor of 5) for a benchmark where the masses of neutralino and chargino are similar and have a value of 200 GeV, the mass of the $\tilde{\tau}$ is 195 GeV and the mass of the LSP is assumed to be zero [20]. No excess above the SM prediction was observed and upper limits were set on the cross sections for pair production of charginos and neutralinos with two associated jets, assuming the stau to be the lightest slepton and with $m(\tilde{\tau}) < m(\tilde{\chi}_1^\pm)$. Figure 2.13 (right) shows the observed limit on production cross sections for the very compressed mass-spectrum scenario where the mass difference of $\tilde{\chi}_2^0/\tilde{\chi}_1^\pm$ and the stau is 5 GeV. Two benchmark scenarios are considered. The results of particular interest are for the case where the mass difference between the $\tilde{\chi}_1^0$ and $\tilde{\chi}_2^0/\tilde{\chi}_1^\pm$ particles was set to 50 GeV. An upper limit of 170 GeV was set for the $\tilde{\chi}_2^0$ and $\tilde{\chi}_1^\pm$ masses. These scenarios with 50 GeV mass gap and light $\tilde{\tau}$ were not excluded by other CMS electroweak SUSY searches (e.g. in the trilepton channels [146]). The VBF topology was shown to be a complementary search tool for the overall electroweak SUSY search program at CMS. The analysis presented in this document uses the VBF topology as a key ingredient to reach to our conclusions as indicated in the analysis chapter.

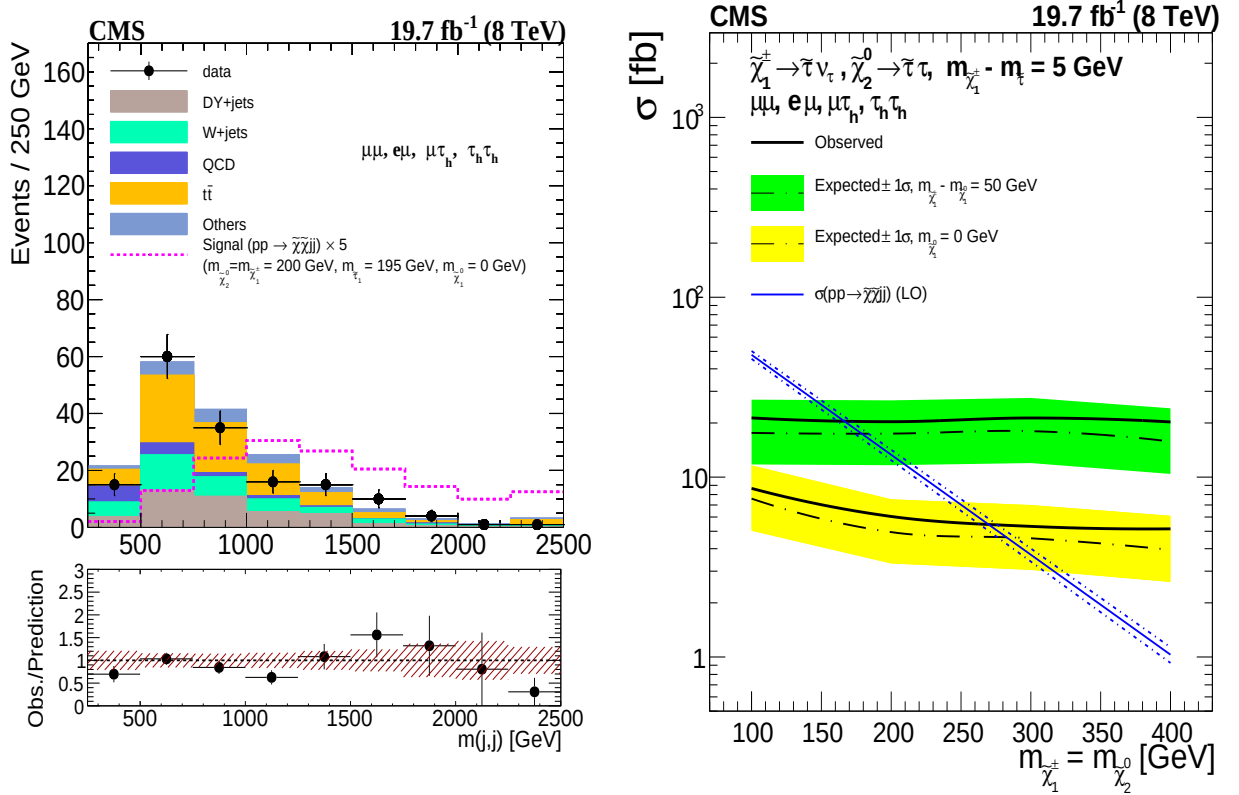


Figure 2.13: (Left) Dijet invariant mass distribution for the combination of all dilepton channels in missing transverse energy + VBF events, comparing data and the SM expectation [20]. (Right) Combined 95% C.L. upper limits on the cross section as a function of $m_{\tilde{\chi}_1^\pm} = m_{\tilde{\chi}_2^0}$ for the scenario where $m_{\tilde{\tau}} = m_{\tilde{\chi}_1^\pm} - 5 \text{ GeV}$ in two different $m_{\tilde{\chi}_1^0}$ definitions: $m_{\tilde{\chi}_1^0} = m_{\tilde{\chi}_1^\pm} - 50 \text{ GeV}$ (compressed mass spectra) and $m_{\tilde{\chi}_1^0} = 0 \text{ GeV}$ (large mass gap) [20].

Chapter 3

The LHC accelerator

3.1 The LHC accelerator complex

The LHC [147] accelerator is located in the French-Swiss border within a tunnel of 27 km of circumference with a mean depth of 100 meters at a slight gradient of 1.4%. The depth varies between 50 m (towards Lake Geneva) and 175 m (under the Jura mountains). The accelerator contains two adjacent parallel beam pipes enclosed with superconducting magnets, where each pipe conducts a beam of particles moving around the ring. The two particle beams traveled in opposite directions with respect to each other and intersect in four different points within the accelerator, where the collisions occur. Each beam is formed by bunches of protons, containing approximately 10^{11} protons. Before injection into the main accelerator, the protons are accelerated by a number of systems that successively increase their energy as it is shown in Figure 3.1. The first system is the linear accelerator LINAC 2. This accelerator allows the protons to reach an energy of 50 MeV. After the LINAC 2, the protons are passed to a synchrotron accelerator, called Proton Synchrotron Booster (PBS) where their energy is increased up to 1.4 GeV. Then, the protons are injected into the Proton Synchrotron (PS) which accelerates them until they reach an energy of 25 GeV. After the PS, the protons are injected into the SPS (Super Proton Synchrotron) which is used to increase their energy up to 450 GeV. Finally, the protons enter the LHC accelerator-collider, where each proton beam is accelerated via eight Radio-Frequency (RF) cavities providing strong electric fields capable of increasing the energy of each proton bunch up to 7 TeV, for a nominal of 14 TeV in the center of mass reference system. The LHC operated at 3.5 TeV energy per proton beam in 2010 and 2011 and at 4 TeV in 2012. In 2015 the LHC increased the proton beam energy to 6.5 TeV [149,150], and has been producing proton-proton collisions at the center of mass energy of 13 TeV since 2015 through 2018. As the energy increases, the increase in speed is minimal since the protons travel very close to the speed of light, c , in the accelerators. Figure 3.2 shows the relationship between energy and speed in terms of units of c . Table 3.1 shows the kinetic energy for the set of accelerators used to inject protons to the LHC. The LHC is also used for colliding heavy ions (p-Pb, Pb-Pb, Xe-Xe). In a regular year of operation, the LHC collides beams of protons during around 6 to 7 months and one month is dedicated to collide heavy ions. In this document, we concentrate on the proton-proton collisions of the LHC.

In order to bend both proton beams on their circular path, 1232 superconducting dipole magnets are installed across the ring, each with 15 metres in length and 35 tonnes of weight. The superconducting magnets are made of niobium-titanium coils that have to be cooled by superfluid helium, at a

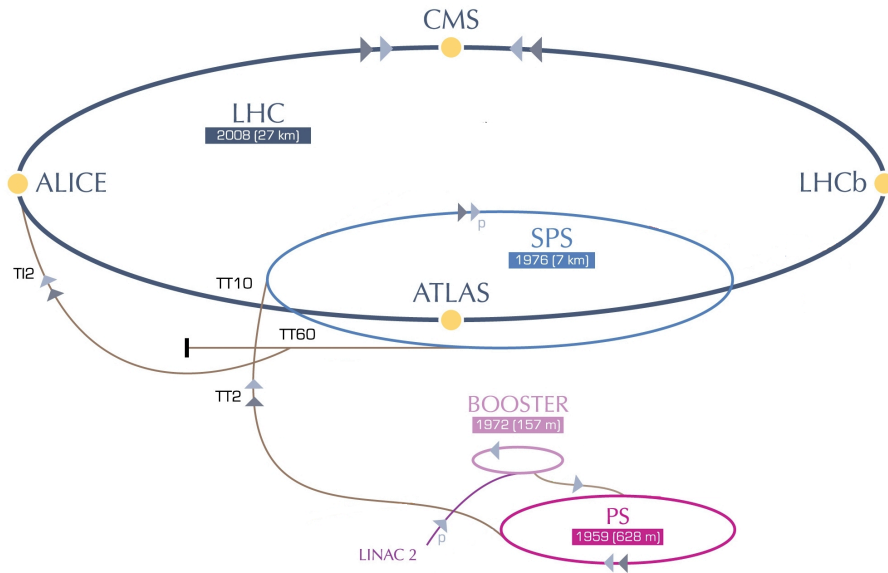


Figure 3.1: Accelerator complex of CERN [148]

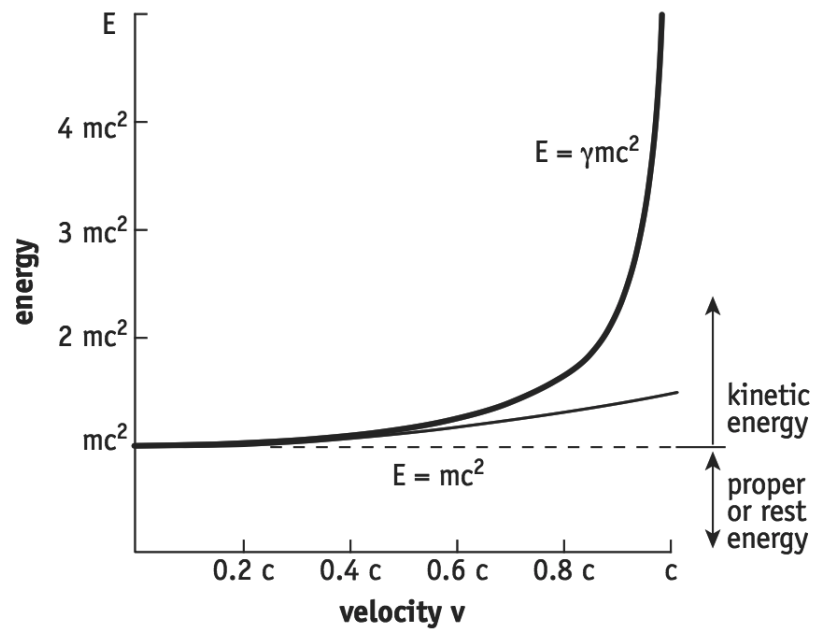


Figure 3.2: Correlation between energy and speed in terms of speed of light units: c [151].

Table 3.1: Kinetic energy achieved in each of the accelerators used in the accelerator complex of CERN [151]

Kinetic energy of a proton (K)	Accelerator
50 MeV	Linac 2
1.4 GeV	PS Booster
25 GeV	PS
450 GeV	SPS
7 TeV	LHC

temperature of 1.9 K, to operate at currents of 11850 amperes to produce magnetic fields of 8.3 Tesla. This high magnetic field is required to bend the 7 TeV proton beams across the 27 km accelerator. About 120 tonnes of liquid helium are required to keep the magnets at this low temperature, 90 tonnes are used in the magnets and the rest in the pipes and refrigerator units [151]. The dipole magnets are also equipped with sextupole, octupole and decapole magnets, which correct for small imperfections in the magnetic field at the extremities of the dipoles [152]. Figure 3.3 shows a cross-sectional area of a dipole magnet with the geometrical distribution of its different elements. Additional 392 quadrupole magnets (each 5-7 metres long) are used to keep the beams focused. Three quadrupoles are used to create the inner triplet magnet which is used to focus the beams to the intersection points where the proton-proton collisions take place. There are eight inner triplets, two of which are symmetrically located at each of the four large LHC detectors, ALICE, ATLAS, CMS and LHCb. Inner triplets reduce the beam transversal dimensions at the interaction points in order to increase the luminosity, making it 12.5 times narrower, from 0.2 millimeters to 16 micrometers in diameter approximately [152].

3.2 Synchrotron radiation

One of the concerns with circular beams is the synchrotron radiation emitted by the accelerated particles when they are bent. The energy loss per unit time (power radiated as synchrotron light) of a charged particle in circular motion is given by [153] :

$$P = \frac{q^4}{6\pi\epsilon_0 m^4 c^5} E^2 B^2 \quad (3.1)$$

Where:

- q is the charge of the particle
- E is the particle energy
- B is the magnetic field bending the particle in circular motion
- m is the particle mass
- c is the speed of light
- ϵ_0 is the permittivity of free space

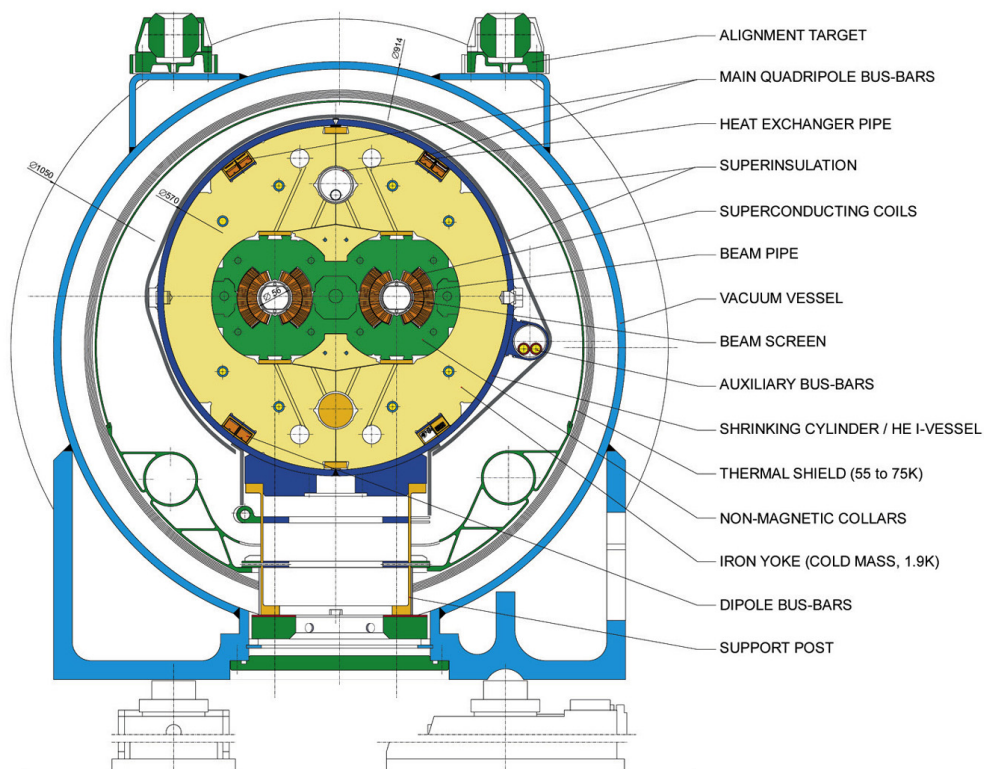


Figure 3.3: Transversal view of an LHC dipole magnet [147].

The energy loss due to synchrotron radiation limited the center of mass energy of LEP (electron positron machine that operated in the same tunnel previously to the construction of the LHC) to $\sqrt{s} = 209$ GeV [154]. Since the mass of the proton is about 2000 times higher than the mass of the electron, protons can be accelerated to much higher energies with minimum effects from synchrotron radiation. Hence, higher proton energies at the LHC can be achieved with the only constraint of achieving very high magnetic fields to bend the beams, which requires very high currents circulating in the superconducting coils of the dipole magnets.

3.3 RF cavities

Radiofrequency (RF) cavities are metallic chambers with an electromagnetic field inside them. The main goal of the cavities is to accelerate charged particles [155]. The charged particles that enter into the cavity increase their energy due to a force produced by the electromagnetic field that push them forward along the accelerator beam. The time of arrival of the charged particles (protons for the LHC) is crucial. Protons that arrive with the right time and energy do not see any accelerating voltage. But protons that arrive later/early are accelerated/decelerated with the overall effect of organizing protons with the same energy in groups called "bunches" [156]. Figure 3.4 shows the energy and time variation of two particles (A and B) over several turns in an RF cavity. Particle A always passes through the cavity when there is no RF voltage (A is called the synchronous particle, as it is exactly synchronised with the RF frequency), but particle B receives an acceleration or deceleration depending if it arrives before or after particle A. Particle B would experience synchrotron oscillations, this means it would oscillate around similar position of particle A under the influence of the RF cavity [155]. As explained in Section 3.1, protons arrive to the LHC ring with an energy of 450 GeV, and then they are accelerated to 7 TeV, completing approximately 11245 turns per second through all the circumference of the accelerator. During this process the bunches experience an acceleration every time they pass through the cavities, the final energy (7 TeV) is accomplished in roughly 20 minutes, this implies that the protons have passed through the cavities more than 1 million times. The total number of cavities in the LHC are 16 (eight cavities per beam). They are stored in four cryomodules (cylindrical refrigerators) that keep the temperature at 4.5 K to assure that the cavities operate in a superconducting state, to avoid energy losses due to electrical resistance [156]. Each cavity delivers a maximum voltage of 2 MV at 400 MHz, and an accelerating field of 5 MV/m [151].

3.4 Luminosity

An important quantity, that measures the ability of a particle accelerator to produce the required number of interactions is called the instantaneous luminosity (usually just called luminosity), which is the proportionality factor between the number of events produced per second of a certain type of process, $\frac{dN}{dt}$, and the cross section σ for that specific process:

$$\frac{dN}{dt} = \sigma \cdot \mathcal{L} \quad (3.2)$$

The units of luminosity are therefore $\text{cm}^{-2}\text{s}^{-1}$. The luminosity is directly proportional to the number of bunches colliding, n_b , the bunch revolution frequency, f , the number of protons per bunch in each

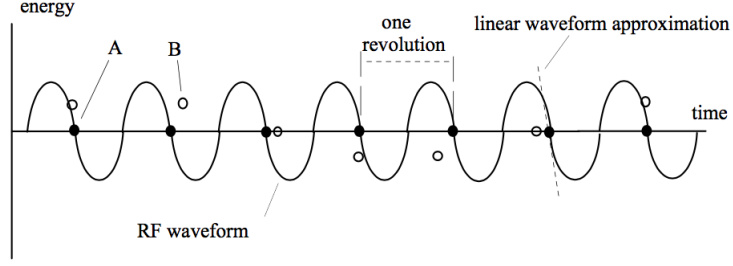


Figure 3.4: Energy and time variation of two particles over several turns in an RF cavity. Particle A is synchronized with the RF frequency while particle B oscillates around similar position as particle A [155].

beam, N_1 and N_2 , and inversely proportional to the transverse area of the beams, according to the equation [157]:

$$\mathcal{L} = \frac{n_b f N_1 N_2}{4\pi\sigma_x\sigma_y} \mathcal{F}, \quad (3.3)$$

where σ_x, σ_y denote the gaussian transverse profiles of the beam at the Interaction Point (IP), $\mathcal{F}$ is a geometric factor that depends on the beam crossing angle at the interaction point, given by the following equation:

$$\mathcal{F} = \frac{1}{\sqrt{1 + \left(\frac{\sigma_s}{\sigma_{\text{xing}}} \frac{\alpha}{2}\right)^2}}, \quad (3.4)$$

where σ_s is the root mean square bunch length, σ_{xing} is the transverse beam size in the crossing plane, and α is the full crossing angle. The integrated luminosity is defined as the instantaneous luminosity integrated over a period of time:

$$L = \int \mathcal{L} dt \quad (3.5)$$

The LHC has been designed to reach a nominal luminosity of $10^{34} \text{cm}^{-2} \text{s}^{-1}$, with a rate of collisions between proton beams of 40 MHz corresponding to crossings of the two beams at the interaction points every 25 ns. With this luminosity, 40 simultaneous interactions on average are expected per bunch crossing [151, 157] which are called pileup events. Also a multiplicity of 100 charged particles is expected per interaction, resulting in nearly 4000 particles detected per each beam crossing. During the first run of the LHC (2011 and 2012), the collider reached, a peak luminosity of $7.7 \times 10^{33} \text{cm}^{-2} \text{s}^{-1}$ and delivered an integrated luminosity of $\sim 25 \text{fb}^{-1}$ to each of the major experiments, ATLAS and CMS [158]. During the second run of the LHC (2015 through 2018), the collider reached a peak luminosity of $\sim 2.1 \times 10^{34} \text{cm}^{-2} \text{s}^{-1}$ and has delivered up to now, to each experiment, an integrated luminosity of $\sim 118.8 \text{fb}^{-1}$ at $\sqrt{s} = 13 \text{TeV}$. Figure 3.5 shows the delivered integrated luminosity to CMS during stable beams for proton-proton collisions since 2010 through 2018 and Figure 3.6 shows the peak and integrated luminosity projection for the LHC until 2035. So far, only $\sim 3\%$ of the total expected LHC integrated luminosity (until 2035) has been recorded.

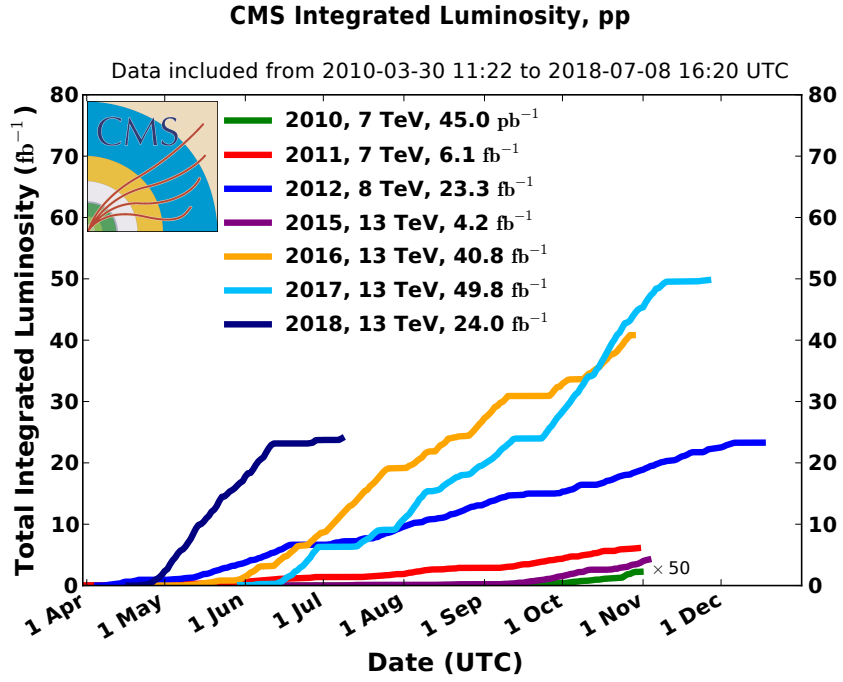


Figure 3.5: Cumulative luminosity versus day delivered to CMS during stable beams for proton-proton collisions. This is shown for 2010 (green, amplified by a factor of 50), 2011 (red), 2012 (blue), 2015 (purple), 2016 (orange) and 2017 (light blue) and 2018 (dark blue) data-taking periods [159].

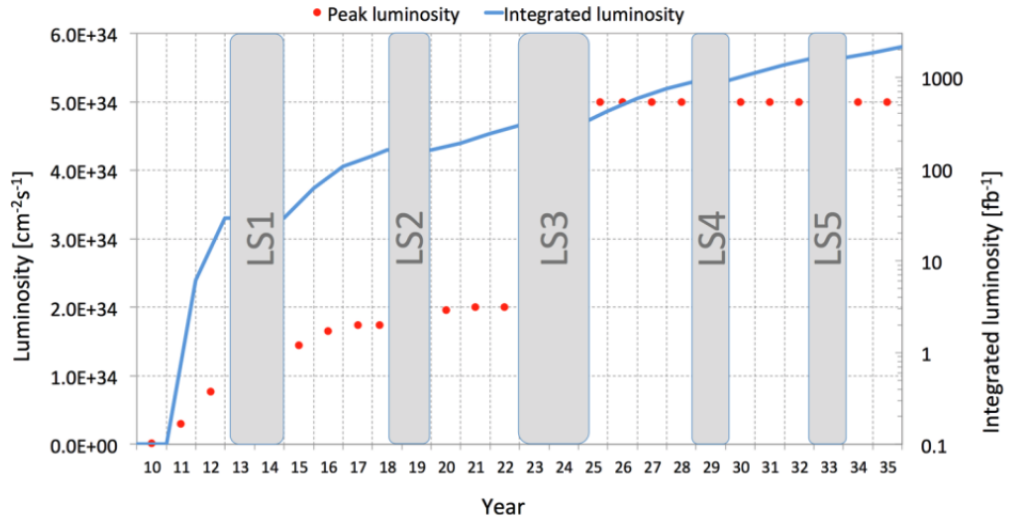


Figure 3.6: Peak and integrated luminosity projection for the LHC until 2035 [158].

3.5 Pileup

As it was already explained in section 3.3, protons inside the LHC travel in groups called bunches (instead of a continuous stream). According to equation 3.3, one way to increase the luminosity is increasing the number of particles in the colliding beams. To achieve that goal, the beam, made of proton bunches, has to be squeezed as narrowly as possible. Pileup is when more than one proton-proton collision takes place every time that the bunches cross one another [160, 161]. The CMS experiment was designed to reconstruct in every 25 ns bunch crossing an average pileup of around 25 interactions per crossing. Nevertheless, current pileup environment is greater than 25, which represents an additional challenge for the CMS tracker. One of the challenges to understand what happens in a single proton-proton collision is to subtract the pileup produced by extra collisions. There are several methods to extract the radiation produced by pile up, that is, the radiation that comes from additional collisions to the primary interaction, which are described in reference. [162]. Pileup depends on instantaneous luminosity. The probability to have n proton-proton interactions in the same bunch crossing follows a poisson distribution:

$$P(n, \mu) = \frac{\mu^n}{n!} e^{-\mu} \quad (3.6)$$

Where μ is the average number of interactions defined by:

$$\mu = \mathcal{L} \sigma_i \Delta t \quad (3.7)$$

$\mathcal{L}$ is the instantaneous luminosity, σ_i is the inelastic cross section and Δt the proton bunch crossing period (25 ns). Figure 3.7 shows the mean number of interactions per bunch crossing for proton-proton collisions at $\sqrt{s} = 13$ TeV measured in the 2016 and 2017 data periods. The average number of interactions was 27 (33) for 2016 (2017). The shape of the distributions depend on the instantaneous luminosities for each proton bunch collision. The two peaks observed for the 2017 data sample is because there were in average two main instantaneous luminosities at which data were taken, as can be observed in Figure 3.8.

3.6 LHC experiments

The LHC has a total of four major experiments distributed around the ring. These experiments (detectors) are located in underground caverns built in the points where the interception of the beams occurs. Two of them are the general purpose detectors ATLAS [163] and CMS [164] which have a broad physics program, mainly to study the Higgs boson and searching for physics beyond the SM. These two detectors were designed to investigate a large variety of physics topics such as SM physics, Higgs searches, SUSY searches, heavy mass resonances, among others. Having two independently designed detectors is vital for confirmation of any new discovery. There are two medium-size experiments, ALICE [165] and LHCb [166] that have specialized detectors. The ALICE experiment is studying the quark gluon plasma produced by collisions of heavy-ions. The LHCb experiment is measuring the parameters of CP violation in the production of b-hadrons. There are two small detectors located near CMS and ATLAS, which are TOTEM [167] and LHCf [168], respectively. The TOTEM detector uses the spray of particles emitted in the forward direction, escaping from the CMS detector, to effectively study diffractive events, measure the proton-proton total cross section and the luminosity delivered by the LHC. Similarly, the LHCf detector is used by

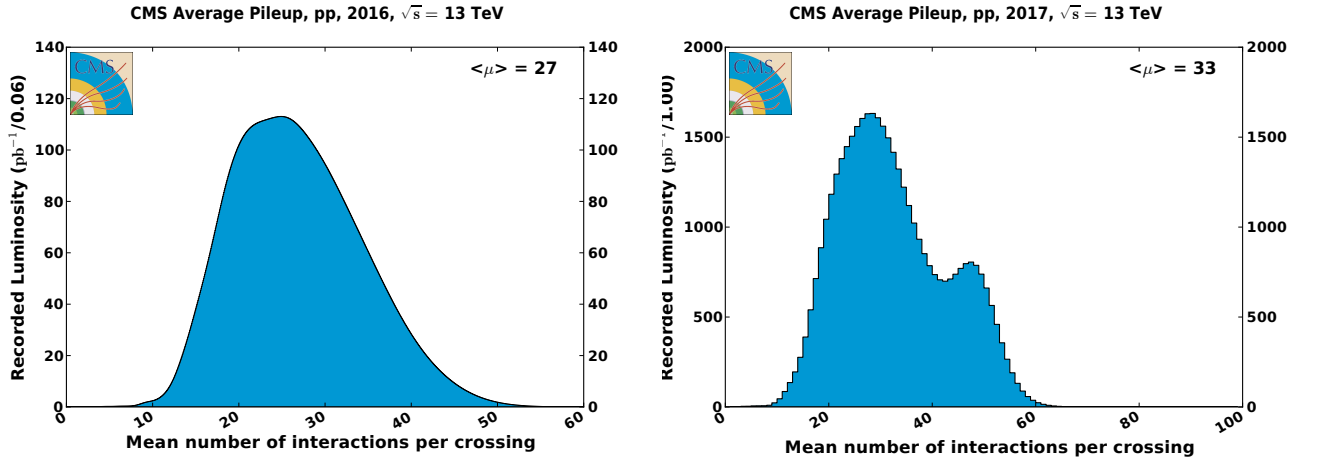


Figure 3.7: Mean number of interactions per bunch crossing for proton-proton collisions taken during 2016 (left) and 2017 (right) at $\sqrt{s} = 13$ TeV [159].

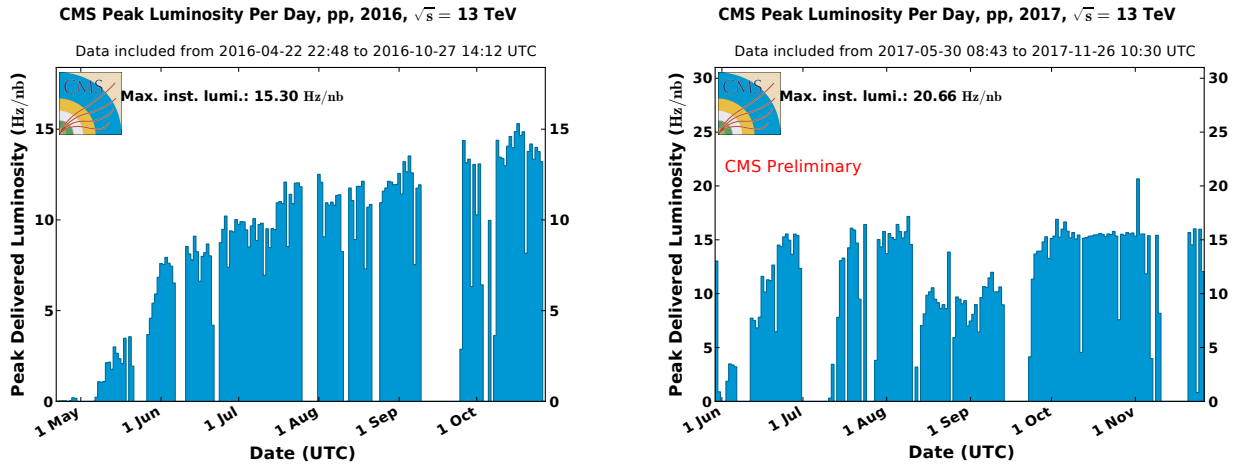


Figure 3.8: Peak instantaneous luminosity on a day-by-day basis for proton-proton collisions taken during 2016 (left) and 2017 (right) at $\sqrt{s} = 13$ TeV [159].

ATLAS to measure the energy and number of neutral pions produced by the LHC for the purpose of explaining the origin of ultra-high-energy cosmic rays.

Table 3.2 lists some of the most relevant parameters of the LHC already mentioned in this chapter.

Table 3.2: Relevant LHC parameters [151].

Quantity	Number
Circumference	26 659 m
Dipole operating temperature	1.9 K (-271.3 °C)
Number of magnets	9593
Number of main dipoles	1232
Number of main quadrupoles	392
Number of RF cavities	8 per beam
Nominal energy, protons	7 TeV
Nominal energy, ions per nucleon	2.76 TeV/u
Peak magnetic dipole field	8.33 T
Min. distance between bunches	~ 7 m
Design luminosity	$10^{34} \text{cm}^{-2} \text{s}^{-1}$
No. of bunches per proton beam	2808
No. of protons per bunch (at start)	1.2×10^{11}
Number of turns per second	11 245

Chapter 4

The CMS Experiment

The CMS detector is one of the two general purpose detectors at the LHC designed to explore the physics in the TeV scale, including the electroweak symmetry breaking mechanism and physics beyond the Standard Model. The detector has a cylindrical geometry with 28.7 meters of length, 15.0 meters in diameter and a weight of approximately 14,000 tonnes. The experiment is located 100 meters underground in Cessy, France near the Swiss border. The CMS detector contains subsystems that are designed to measure the energy and the momentum of photons, electrons, muons and other products of the collisions. Figure 4.1 shows a diagram of a transverse section of the CMS detector with the different subsystems of the experiment and the type of particles that can be detected by each subsystem. The innermost layer detector is made of silicon pixels. This detector is surrounded by the electromagnetic calorimeter which is made of scintillating crystals, which in turn is surrounded by a sampling calorimeter for identifying hadrons (the hadronic calorimeter). The silicon detector and the calorimeters are compact enough to be within the CMS solenoid which generates a magnetic field of 3.8 Teslas. Around the Solenoid there are the muon detectors which are located within an iron yoke structure used for the return of the magnetic field. The proton beam collisions occur in the center of the detector. Nevertheless, the collision probability is very small given the measurements performed at 8 TeV and 13 TeV. The measured inelastic cross section for proton-proton collisions at a centre-of-mass energy of $\sqrt{s} = 13$ TeV by the CMS experiment is 71.3 mb [170]. Figure 4.2 shows a diagram of the CMS experiment and its different subdetectors.

Coordinate system

Figure 4.3 shows a scheme of the coordinate system used in the CMS experiment. The experiment uses a right-handed coordinate system with the origin located at the point where the beam intersection takes place (nominal interaction point). The x-axis points towards the center of the LHC ring, the y-axis points up, and the z-axis points towards the direction of the beam (counterclockwise). The polar angle (θ) is measured from the z-axis, and the azimuthal angle (ϕ) is measured relative to the x-axis in the x-y plane. Instead of the θ angle, the pseudorapidity (η) is used. Defined as follows:

$$\eta = -\ln\left(\tan\frac{\theta}{2}\right) \quad (4.1)$$

The use of pseudorapidity instead of the polar angle is motivated because the difference of pseudorapidity between two particles is invariant under Lorentz boosts in the beam direction [171]. The

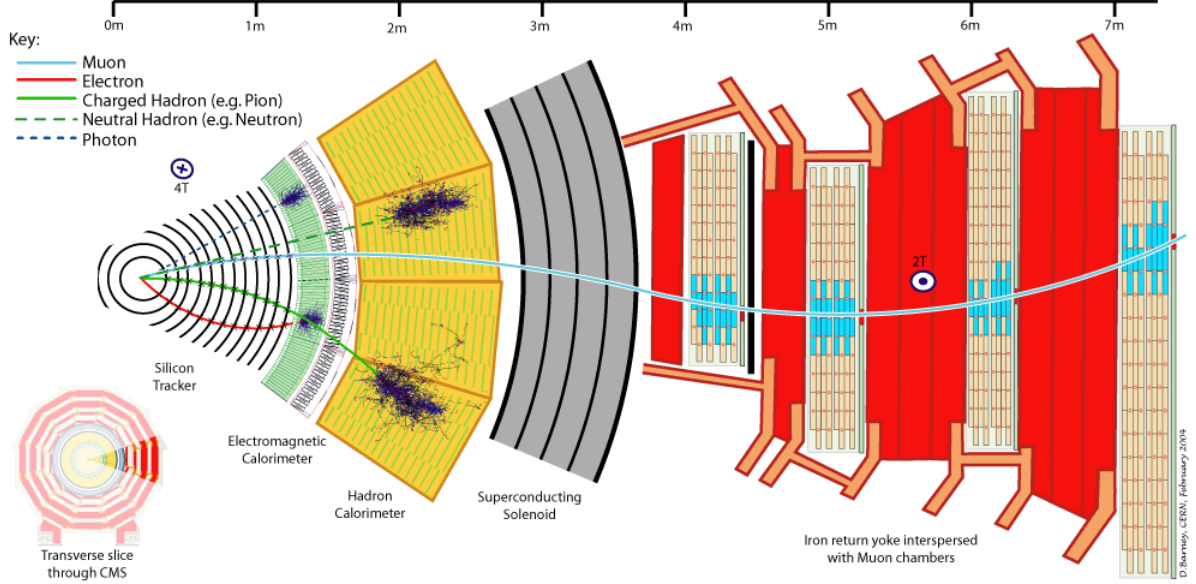


Figure 4.1: Transverse Section of CMS detector showing the subsystems that identify different types of particles [169].

Lorentz-invariant distance between two relativistic objects is defined as follows:

$$\Delta R = \sqrt{(\Delta\eta)^2 + (\Delta\phi)^2}. \quad (4.2)$$

One of the variables mostly used in the experiment is the transverse momentum defined as:

$$p_T = \sqrt{p_x^2 + p_y^2} \quad (4.3)$$

where p_x and p_y are the projection of the momentum vector on the x-axis and y-axis. The momentum conservation in the transverse plane is commonly used to obtain more information about the final states because the initial momentum (before the proton-proton collisions) in the transverse plane is zero.

4.0.1 Tracker

The region of the detector that is closest to the beam is the tracker. Its function is to record the traces left by the charged particles in their path. These traces are used to reconstruct the trajectory of these particles with high precision and also to reconstruct secondary vertices of long-lived particle decays. The tracker is located on a cylinder of 2.5 m in diameter and 5.8 m long around the point of collision. It must have a high granularity due to the high multiplicity of particles by all the collisions (including pileup). It must also allow associating the identified particles with their respective bunch crossings. A transversal section of the tracker with its subsystems is shown in figure 4.4. The tracker is formed by the pixel detector that is in the center, and by the silicon strip detectors that are around the pixel.

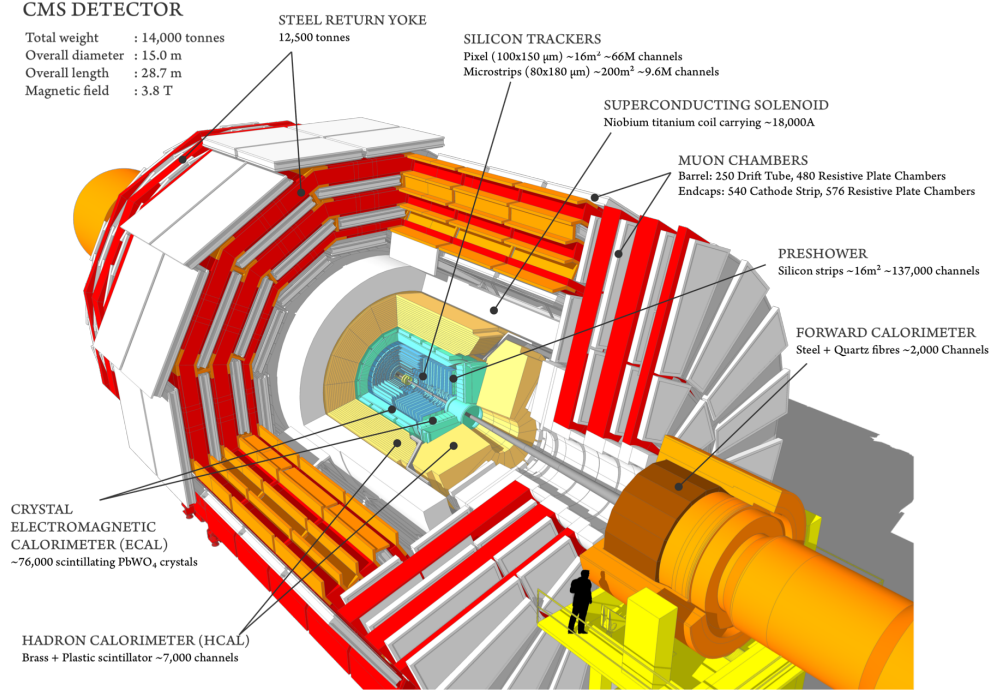


Figure 4.2: CMS experiment layout. The central region of the detector is called the barrel, which is composed of five wheels. The caps in the two sides of the experiment are called the endcap region. [169]

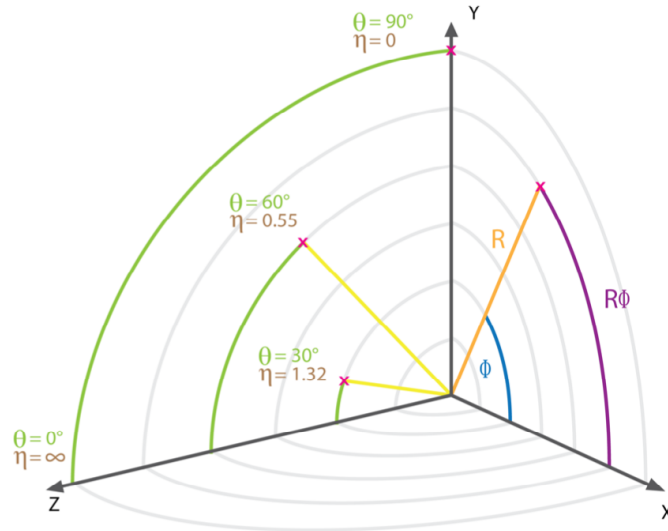


Figure 4.3: Coordinate system used in the CMS experiment. The diagram shows the definition of pseudorapidity (η) and azimuthal angle (ϕ) [172].

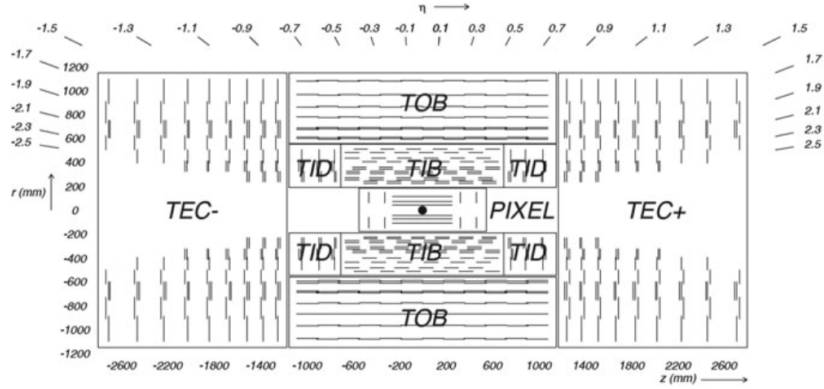


Figure 4.4: Sectional view of the tracker [173].

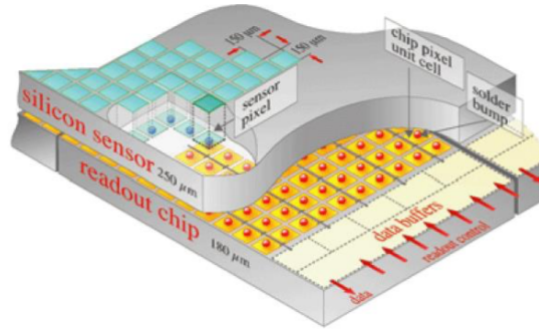


Figure 4.5: Schematic representation of the CMS silicon pixel detector.

Tracker Subsystems

Pixel

The most internal detector of the tracker is the pixel detector that is installed within a distance of 20 cm from the beam pipe. It has an active area of 1.1 m^2 . The pixel consists of four cylindrical barrel layers located at 3, 6.8, 10.2 and 16 cm from the interaction point and 6 endcap layers, 3 on each side of the barrel layers, located at 24.1, 34.5 and 45.5 cm from the interaction point. Figure 4.5 shows a schematic depiction of the pixel silicon detectors installed in CMS. There are a total of 66 million pixel sensors installed with a size of $100 \mu\text{m}$ by $150 \mu\text{m}$. The sensors from the barrel layers are oriented with the smaller pitch pointing in the azimuthal direction (ϕ), while the sensors from the endcap layers are oriented in the radial direction. The entire detector provides a coverage in $|\eta|$ from 0 to 2.4. This detector is used to reconstruct primary and secondary vertices from unstable particles that decay within the detector volume (e.g. b-quarks), which are used later for the identification of b-jets.

Silicone Strips

The CMS silicone tracker has 9.3 million of micro-strips with an active area of 200 m². It surrounds the pixel detector and covers a distance up to 116 cm from the beam pipe. The entire detector provides a coverage in pseudorapidity of $|\eta| < 2.4$. The silicon micro-strip detector is composed of three subsystems:

- The Tracker Inner Barrel (TIB) and Tracker Inner Disks (TID) composed of 320 μm thick silicon sensors. The TIB is parallel to the z-axis and is composed of 4 layers. It is located in a radial region between 20 and 55 cm. The strips are oriented parallel to the z-axis (beam axis) with a strip pitch that varies between 80 μm and 120 μm . The TID is composed of three disks located between 80 and 90 cm on the z-axis with a pseudorapidity coverage of $|\eta| < 2.5$. The strips are oriented radially with a strip pitch that varies between 100 μm to 141 μm .
- The Tracker Outer Barrel (TOB) surrounds the TIB and TID. It is also parallel to the z-axis and is formed by six layers. It has a length of 2.18 meters. Their sensors have a thickness of 500 μm with a strip pitch of 122 μm for the first four layers and 183 μm for the last two.
- The Tracker Endcap (TEC) is composed of 9 disks perpendicular to the z-axis, each of them divided into 16 angular subsections. This region is located between 134 and 282 cm along the z-axis and provides a pseudorapidity coverage of $|\eta| < 2.5$. Their sensors have a thickness of 320 μm for the first four rings and 500 μm for the others. The average strip pitch varies from 97 μm to 184 μm .

Tracker Resolution

The momentum resolution (measured with muons) for a charged particle between 1 and 10 GeV is less than 1 % for values of $|\eta| < 1.9$. This value changes to less than 2 % for values of $|\eta| \geq 1.9$. For 100 GeV particles, this resolution is 2 % for $|\eta| < 1.6$ and is degraded to 7 % for $|\eta|$ values up to 2.4. The tracker traces are also used to reconstruct secondary vertices of long-lived particle decays (the resolution for the impact parameter in the tracker is 15 μm). The charged hadrons are reconstructed in the tracker with an efficiency of 95 % for p_T greater than 10 GeV. In this thesis, the tracker is fundamental for the reconstruction of muons, since the muon reconstruction algorithms use the traces left in the tracker to make matching with the traces left in the muon chambers to determine the whole trajectory of the muons along all the detector. Likewise, it is fundamental for the electron reconstruction (which later deposit all their energy in the electromagnetic calorimeter). Due to the decay of the tau in charged hadrons, a proper identification of the hadrons in the tracker allows a good reconstruction of tau. The identification of b-jets depends on the impact parameter and secondary-vertex resolution of the tracker. A good reconstruction of b-jets in the tracker is fundamental in our analysis to discriminate the VBF jets from b-jets.

4.0.2 Electromagnetic Calorimeter

The electromagnetic calorimeter (ECAL) is designed to detect and absorb the energy deposited by photons and electrons. The interaction between these particles and the material produces an electromagnetic shower formed by cascades of electrons and photons (a representation of the shower is shown in Figure 4.6). The cascade is triggered when an electron or photon enters into a medium

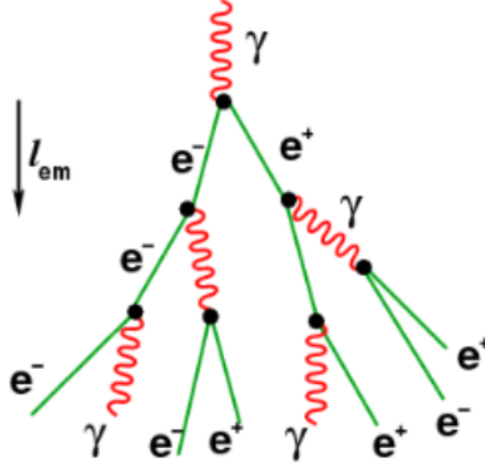


Figure 4.6: Schematic view of the development of an electromagnetic shower [174].

of high density (in this case the calorimeter) and begins to lose energy due to the interactions with the medium. Photons lose their energy due to electron-positron conversions while electrons lose their energy due to bremsstrahlung radiation. This process continues until the loss of energy can no longer produce secondary particles. The longitudinal and transverse behaviour of electromagnetic shower allows determining the energy of the particle that generates it. This is measured with the radiation length (X_0) and the Moliere radius. The Moliere radius is the radius in which 90 % of the electromagnetic shower is deposited. Radiation length is the distance an electron or a photon travels until its energy is reduced by a factor of $1/e$. For this reason, the electromagnetic calorimeter was designed with sufficient width in terms of radiation lengths, so that all the photons and electrons, independently of their initial energy, deposit their total energy in the material. Additionally, the ECAL was designed with a material that has a small Moliere radius to allow good granularity.

The ECAL was built with lead tungstenate crystals ($PbWO_4$ see Figure 4.7). This material has a radiation length of $X_0 = 0.89$ cm and small Moliere radius (2.2 cm). The crystals have an area of 22 mm x 22 mm and a depth of 230 mm in the barrel region; and 3 cm x 3 cm x 22 cm in the endcap region. All the crystals are located in a set of carbon fiber matrices to keep them optically isolated. In the barrel region there are 61200 crystals organized in 36 modules; each module contains 1700 crystals. In the endcaps there are 7324 crystals, in this region, there is also the Preshower detector that allows distinguishing between high energy isolated photons and low energy photons. All the crystals of the ECAL provide a coverage in pseudorapidity of $|\eta| < 3.0$. The arrangement in which the crystals are distributed in the electromagnetic calorimeter is shown Figure 4.8. When electrons and photons cross the crystals, they emit a pulse of light because the crystals are made of a scintillating material. This pulse is then detected by some photodetectors that convert the pulse of light into an electrical signal. The design of the electromagnetic calorimeter was directed towards the requirements imposed for the search of the Higgs boson in the channel $H \rightarrow \gamma\gamma$ where a peak should be distinguished in the invariant mass of the two photons with respect to the background.



Figure 4.7: CMS ECAL Lead Tungstate Crystals ($PbWO_4$).

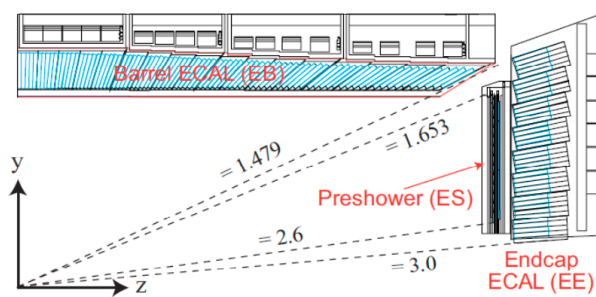


Figure 4.8: Arrangement of the ECAL components.

Then good resolution and fine granularity are required; both things allow to improve the resolution of the invariant mass of the two photons, respectively improving the energy and the measure of the angle of the photons. The fine granularity also allows good identification of $\pi^0 \rightarrow \gamma\gamma$. This calorimeter is inside the solenoid to avoid worsening the resolution in energy, therefore, it has to be compact. The energy resolution for the electromagnetic calorimeter is obtained by making Gaussian fits to the reconstructed energy. This resolution can be parameterized as a function of the energy following the following equation:

$$\left(\frac{\sigma_E}{E}\right)^2 = \left(\frac{S}{\sqrt{E}}\right)^2 + \left(\frac{N}{E}\right)^2 + (C)^2 \quad (4.4)$$

Where S represents statistical fluctuations in the number of secondary particles produced, N represents electronic noise, and C is a constant for calibration errors. The measured values are: $S = 0.028 \sqrt{\text{GeV}}$, $N = 0.12 \text{ GeV}$ and $C = 0.003$, when there is no magnetic field, and without any other material between the beam and the calorimeter. Finally, for photons with transverse energy greater than 100 GeV, the energy resolution of the ECAL is of the order of 0.5 %.

4.0.3 Hadronic Calorimeter

The purpose of the hadronic calorimeter (HCAL) is to stop and measure the energy of the hadrons produced in each event. This calorimeter plays an essential role in the identification of quarks, gluons, and neutrons by measuring the energy and the direction of the jets produced in each event. Additionally, it allows to measure the E_T^{miss} in the events. The E_T^{miss} would be a crucial signal of SUSY particles (like the LSP). For a good resolution in E_T^{miss} , the calorimeter must have a hermetic coverage (that means $|\eta| \leq 5$). The HCAL also allows the identification of electrons, photons, and muons when working in conjunction with the tracker, the electromagnetic calorimeter and the muon systems. The hadronic calorimeter consists of plates of dense material (brass or steel) interspersed with plastic scintillators. When a hadron reaches the layers of dense material, it produces a cascade of particles, called hadronic shower, consisting mainly of pions originating from the strong interaction between the hadron and the atomic nucleus of the material. This cascade can also acquire an electromagnetic component due to the decay of the neutral pion into photons. The particles produced in the cascade are detected by plastic scintillators, which in turn emit pulses of light that are collected by optical fibers and converted into electrical signals. The amount of energy of the incident hadron is measured by the amount of energy deposited in the scintillator material. A hadronic shower is described by the absorption length, which is defined as the average distance traveled by a hadron through a medium before it interacts with the material:

$$\lambda = \frac{A}{N_A \sigma_{\text{abs}}} \quad (4.5)$$

where A is the nuclei weight, N_A is the Avogadro's number and σ_{abs} is the absorption cross-section. In high pseudorapidity regions ($3 < |\eta| < 5$), the Hadronic Forward detector is used, which is located 11 meters on each side of the interaction point, as shown in Figure 4.9. This detector uses a slightly different technology than the one used in the barrel region; it employs steel absorbers and quartz fibers to obtain the signal output. This is because there is much more particle congestion in this region (forward) than in the barrel region. A good resolution and efficiency in this region of η is fundamental for the VBF jets reconstruction, which are required to be within $|\eta| \leq 5$. The energy

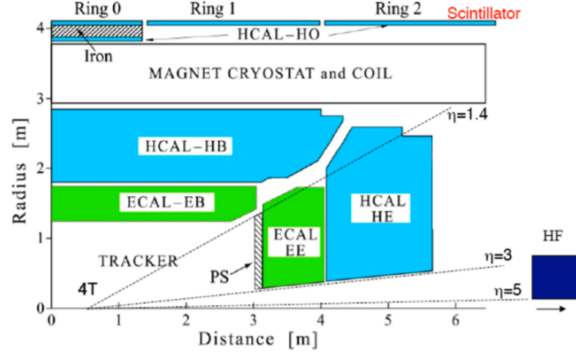


Figure 4.9: Disposition of the HCAL components..

resolution for a pion in the HCAL [173] is given by

$$\left(\frac{\sigma_E}{E}\right)^2 = \left(\frac{138\%}{\sqrt{E}}\right)^2 + (13\%)^2 \quad (4.6)$$

It can be seen that the energy resolution obtained with the HCAL is worse in comparison with the energy resolution obtained with the ECAL. However, this is not a problem because all the information provided by the HCAL is combined with the information coming from the other CMS subdetectors.

4.0.4 Superconducting Solenoid

The solenoid magnet is surrounding the hadron calorimeter. It is 12.9 meters long with an inner bore of 5.9 meters. It is made of niobium-titanium superconducting coils with 2168 turns through the entire solenoid. The coils are refrigerated with liquid helium at a temperature of 4.65 K [175] to allow a current of 19.5 kA that produces a 3.8 Tesla magnetic field. The magnet was initially designed to operate at 4 Tesla, however, currently the field produced is 3.8 Tesla to maximize the magnet longevity. The magnet stores a total energy of 2.7 GJ. The high magnetic field allows the determination of the charge/mass ratio of particles due to the curved track followed by the charged particles that traverse the magnetic field (see Figure 4.1). The external magnetic field is achieved by means of the solenoid holder (the yoke). The yoke is made of an iron frame that diverts the magnetic field lines generated inside the solenoid. Then, outside the solenoid an opposing magnetic field is produced (see Figure 4.10). The yoke and the solenoid add up to a total of 12500 tons [176]. The external magnetic field is important for the detection and identification of muons due to the curve followed by them in the presence of the field. The curved path of the muons will influence the signal obtained by the muon chambers. Afterwards these muon chamber will allow the reconstruction of the muon trajectory (see Figure 4.1). The momentum of a charged particle in the presence of a magnetic field can be written as:

$$p = \gamma mv = qBr \quad (4.7)$$

Where

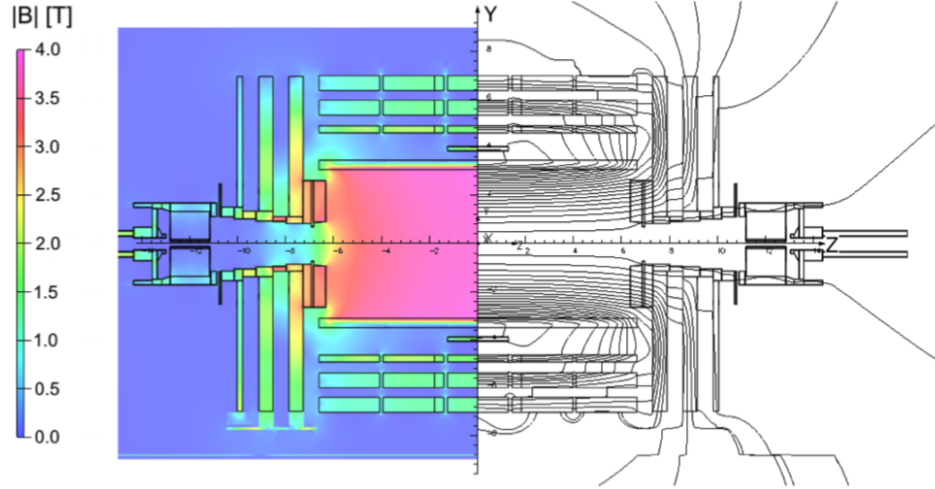


Figure 4.10: Simulation of a 3.8 Tesla magnetic field lines in a longitudinal section of the CMS experiment. The left side shows the absolute value of the field. The right side shows the field lines. Each field line in the drawing represents an increase in the magnetic flux of 6 Wb [178].

- γ is the Lorentz factor: $\gamma = \frac{1}{\sqrt{1 - \frac{v^2}{c^2}}}$
- m is the mass of the particle
- v is the speed of the particle
- q is the charge of the particle
- B is the magnetic field strength
- r is the bending radius

A huge magnetic field is required to measure the momentum of high-energy charged particles accurately. This implies that the momentum resolution depends on the magnetic field, and also on the spatial resolution of the tracker detector. The strong magnetic field together with the excellent tracker resolution allows a 1% muon momentum resolution, for muon momentum up to 100 GeV [177].

4.0.5 Muon System

One of the most important tasks of the CMS experiment is to detect muons. Their presence in the final state of a proton-proton collision is usually a sign of interesting physics. For example, one of the cleanest signatures of the Higgs Boson is its decay with four muons in the final state, also, it is expected that muons will be produced in the decay of a number of potential new particles [179]. The muon system of the CMS experiment has three different type of gaseous detectors. For an accurate measurement of the muon trajectory, in the central region of the barrel ($|\eta| < 1.2$), Drift Tubes (DT)

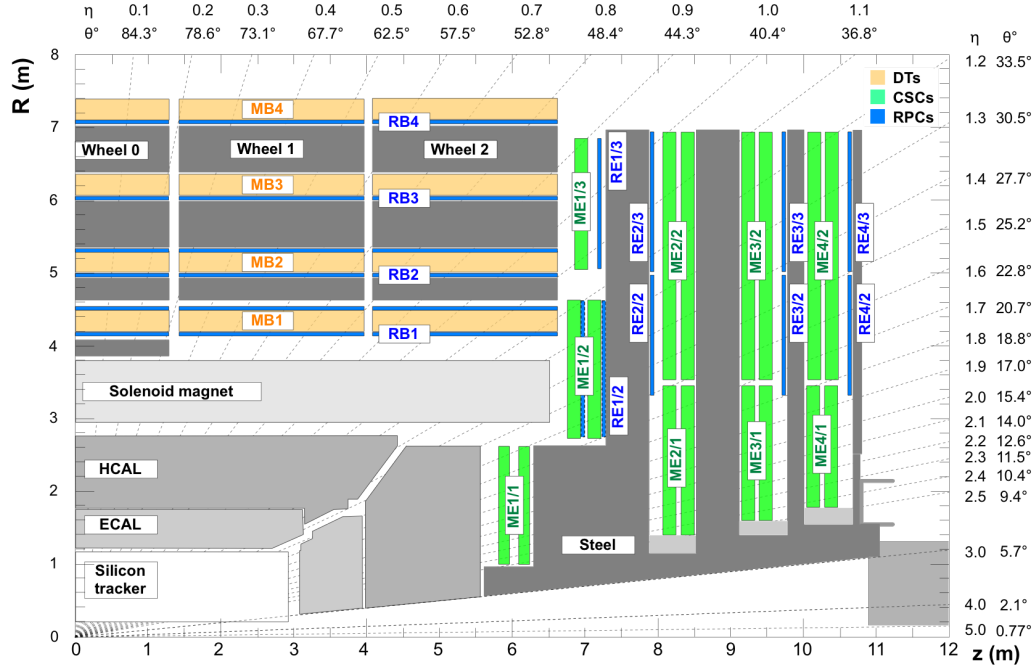


Figure 4.11: Cross section of a quadrant of the CMS detector [181]

are used. Cathode Strip Chambers (CSC), are used in the endcap region ($0.9 < |\eta| < 2.4$) for the same purpose. In both, the barrel and the endcap region ($|\eta| < 1.6$ for L1 trigger and $|\eta| < 1.8$ for reconstruction), Resistive Plate Chambers (RPC) are installed as it is shown in Figure 4.11. RPCs provide a fast signal when there is a muon passing through them, this signal is used to activate the DTs, CSCs and the overall data acquisition system of the CMS experiment [180].

Drift Tubes

Drift tubes (DT) are located in the barrel of the CMS detector (occupying a pseudorapidity region of $|\eta| < 1.2$), where the magnetic field is guided and trapped almost entirely by the iron plates of the yoke. Each tube contains a wire with a large spacing (4 cm), and is arranged in layers. Only the signals in the wires are recorded, resulting in a moderate number of electronic channels used to read the signals from the tubes. An ionizing particle that passes through the tube releases electrons that move along the field lines until reaching the wire that is held at a positive potential.

The coordinate in the plane perpendicular to the wire is obtained with very high accuracy by calculating the time spent by the electrons produced by ionization to migrate to the wire. This time (which is measured with an accuracy of 1 ns) multiplied by the electron drift velocity in the gas results in the distance from the wire. A drift tube chamber occupies a size of $2\text{m} \times 2.5\text{m}$ (see Figure 4.12). The largest drift chambers used in the CMS experiment have a size of $4\text{m} \times 2.5\text{m}$. Figure 4.13 shows the internal structure of a drift tube.

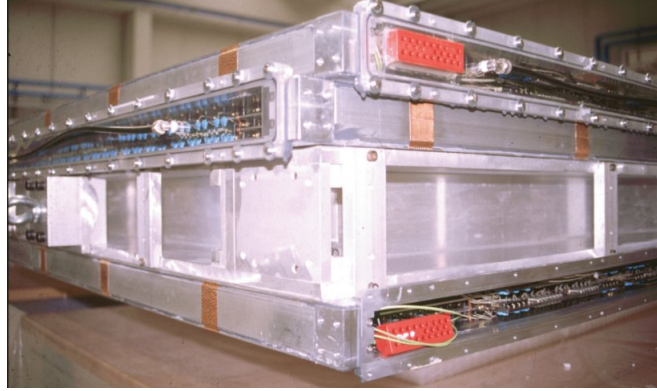


Figure 4.12: Final prototype of a drift tube chamber [182].

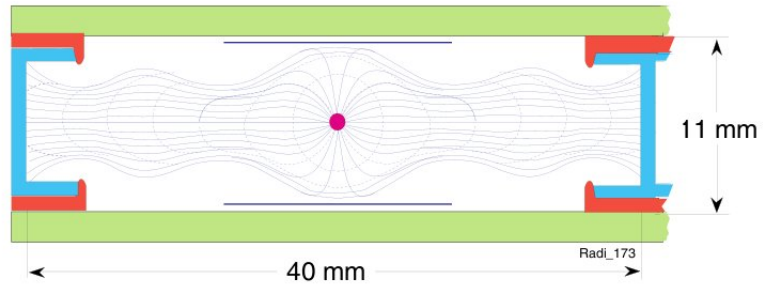


Figure 4.13: Section of a cell in a drift tube showing the drift lines and isochrones. The electrodes are in dark blue and the cathode in light blue color. The anode of the wire is in the center. The applied voltages are 3600 V for wires, 1800 V for the electrodes and -1200 V for cathodes. [182].

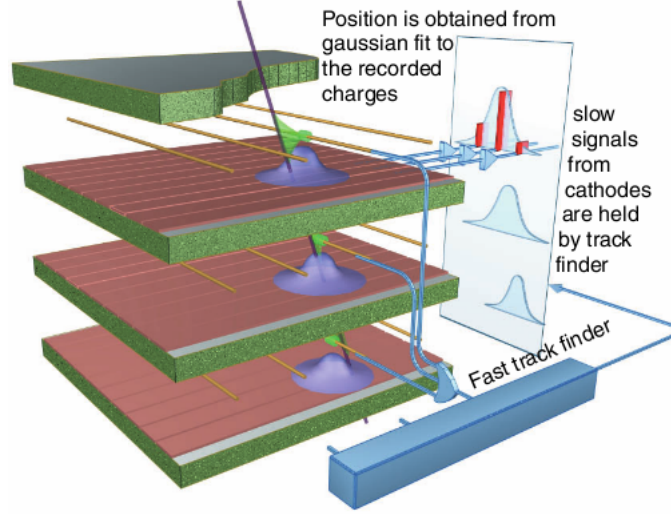


Figure 4.14: CSC chamber and its response to the passage of a charged particle [182].

Cathode Strip Chambers

The cathode strip chambers (CSC) are located in the endcaps of the CMS detector (occupying a pseudorapidity region of $0.9 < |\eta| < 2.4$) where the magnetic field is very intense and very inhomogeneous. They are multiwire proportional chambers where the plane of the cathode is segmented into strips along the wires. A developing avalanche in one wire induces a charge in some of the strips from the plane of the cathode. A CSC is equipped with two coordinates per plane to make detection. This is because of the simultaneous and independent detection of the same trace induced in the wire and in the signal strips. The wires provide the radial coordinate, while the strips provide the angular coordinate ϕ . A diagram of operation of a CSC is shown in Figure 4.14.

Resistive Plate Chambers

The RPCs are gaseous detectors [183] with good time resolution (~ 1 ns with a 2 mm single-gap, 20 ps for 0.1 mm multiple-gap RPCs) and spatial resolution (≤ 10 mm) [184] that can identify the passage of charged particles. An RPC is made of two parallel plates with a phenolic high pressure laminate (bakelite). Resistive plates of resistivity ranging from 10^{10} to 10^{11} Ωcm are separated by a 2 millimeter gas gap as indicated in Figure 4.15. The outer layers of the plates are covered with graphite paint to form the electrodes of the chamber. One electrode acts as the ground, and the other has a specific electric potential for the measurement of the ionization produced in the chamber [185]. Copper strips that are separated by a layer of insulating material (mylar or PET film) receive the output signal (see figure 4.15).

The CMS experiment uses double gap RPC chambers (see figures 4.15 and 4.16). They are operated in avalanche mode with a gas mixture of three components (with 35%-40% of Humidity):

- 95.2% $\text{C}_2\text{H}_2\text{F}_4$ (Tetrafluoroethane, R-134a)
- 4.5% iC_4H_{10} (Isobutane)

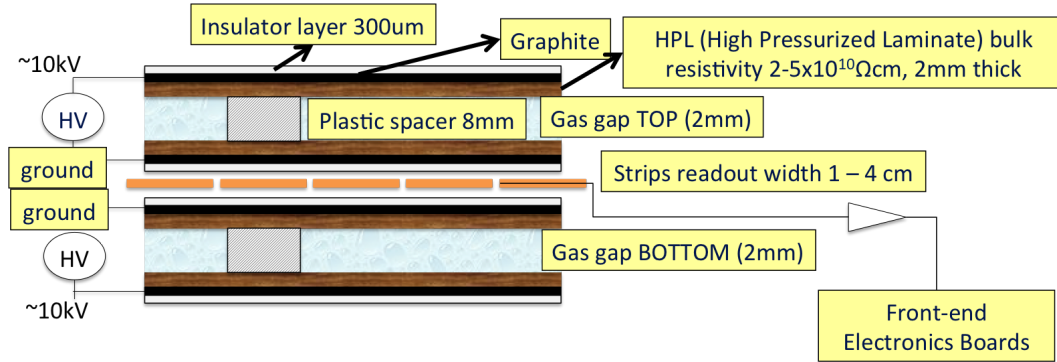


Figure 4.15: Cross-section of a Resistive Plate Chamber [186].

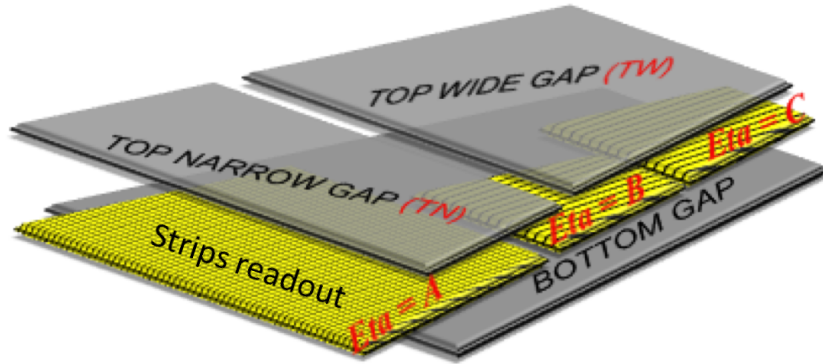


Figure 4.16: Representation of an RPC Double gap chamber in the CMS Experiment [186].

- 0.3% SF₆

Operation Modes

RPCs can operate in different modes. In “streamer mode”, an RPC operates with a strong electric field that produces localized gas discharges in the region near the passage of the ionizing particle, however this mode allows few counts per unit area ($\sim 100 \text{ Hz/cm}^2$). A second mode of operation is the “avalanche mode”, where the electric field between the plates decreases (in comparison with the streamer mode), and the output signal is amplified. The reduction in voltage decreases the charge generated in the ionization and increases the counting capacity of the chamber.

The electrons produced by the ionisation due to the presence of charged particles in the gas are multiplied in avalanche mode by a uniform electric field of about 4.5 kV/mm (see figure 4.17). Typical values of induced charge when the RPC is in avalanche mode are around a few pC. This value could increase by a factor of 50 when the chamber is in streamer mode. The streamer mode is used mainly in cosmic rays, while the avalanche mode is used in applications requiring detection of high fluxes of particles [187]. Once the signal is obtained, it is read out using a capacitive coupling through copper strips in the outer part of the electrodes (See Figure 4.17).

The RPC chambers have a simple mechanical structure, no wires, and are therefore easy to man-

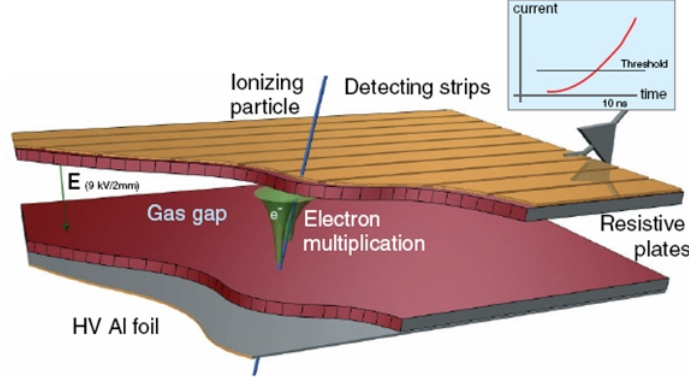


Figure 4.17: Section of an RPC chamber formed by two plates separated by a gap of gas 2mm wide. An electric field of approximately 4.5 kV/mm is applied to the electrodes of the chamber [187].

ufacture. Two RPCs can be used to form a dual chamber structure with common copper strips in the middle (see Figure 4.16). An RPC chamber has the capacity to differentiate ionizations of less than 25 ns between two successive crossings of particle bunches. A muon detector based on RPC chambers can unambiguously identify the beam crossing that originated the muon footprint, even when the system has a high incidence of muons. All these functions can be achieved in high radiation environments with a rate capability of 10^3 Hz/cm^2 .

RPC Chamber Distribution in the CMS Experiment

The barrel region of the CMS experiment is equipped with 480 RPC chambers. Each endcap region (positive and negative) is equipped with 216 chambers. During the long shutdown of the LHC in 2013 and 2014, additional 144 RPC chambers were installed in the endcap region, corresponding to the RE4/2 and RE4/3 stations indicated in Figure 4.11. There are in total 1056 RPC chambers for the entire CMS experiment as shown in Figure 4.18. Given the increased redundancy with these upgrades, higher efficiency is achieved in the Level 1 muon trigger. The extra RPC stations allow the possibility to tighten the trigger quality requirements, producing a reduction on the rate, while keeping a reasonably low threshold with a high instantaneous luminosity [186].

Operation Principle of an RPC

The resistivity of the electrodes mainly determines the rate of counts in the RPC, while the space of the gas determines the time performance of the chamber. Other important parameters for achieving the best performance are the gas density in the chamber and the thickness of the electrodes.

When a charged particle passes through an RPC, it ionizes the gas that is located between the plates, leading to an accumulation of electrons (n_0). Which in turn produces an avalanche of electrons due to the electric field between the plates. This avalanche creates an electrical charge Q_e in the gas with thickness d . The travel of this generated charge towards the anode rapidly induces in one of the electrodes the charge q_e , which represents the signal received by the RPC due to ionization. The electrical source has to move the charge q_s in the circuit outside the gas to compensate the charge collected in the electrodes.

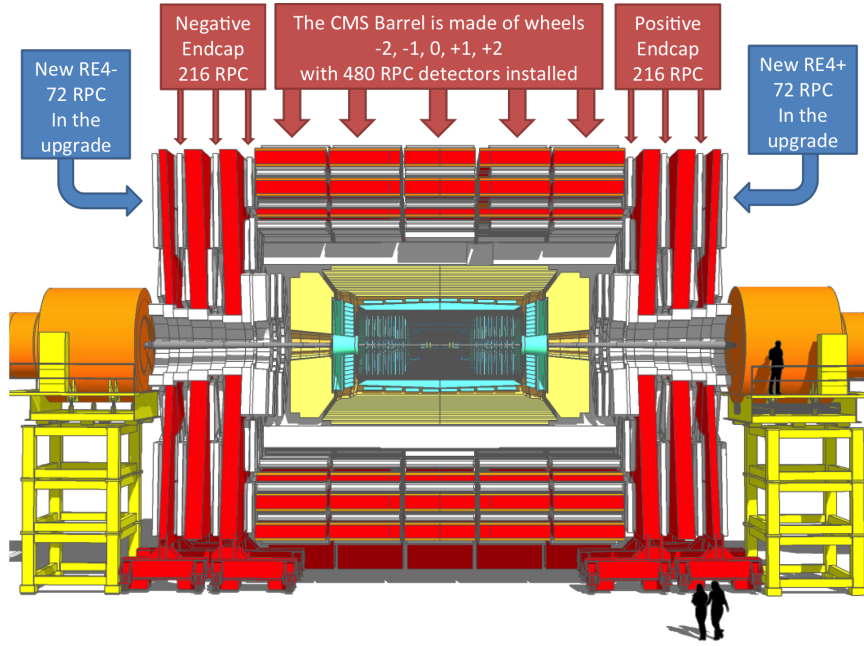


Figure 4.18: RPC Chamber Distribution in the CMS Experiment. [186]

If α is the number of ionizing events per unit length experienced by an electron and the coupling coefficient β is the number of coupling events per unit length, then the effective ionization can be defined as $\eta = \alpha - \beta$. An RPC is in avalanche mode when the condition $\eta d < 20$ is satisfied, in this case, as shown in [185], the average charge induced in one of the electrodes (q_e) is obtained as follows:

$$\langle q_e \rangle = \frac{k}{\eta d} \langle Q_e(d) \rangle = q_{el} n_0 e^{\eta d} \frac{k \lambda}{\eta d (\eta + \lambda)} \quad (4.8)$$

Where:

- $k = \frac{\epsilon_r d / s}{\epsilon_r d / s + 2}$ is a constant depending on the material parameters
- q_{el} is the electron charge
- n_0 is the average size of the charge cumulus where the avalanche was originated
- λ is the cumulus density in the gas (the number of primary cumulus per unit length produced by an ionizing particle)
- ϵ_r is the dielectric constant of the electrode
- d is the thickness of the gas (2 mm in this case)
- s is the thickness of the electrode

For a given value of ηd , the factors k and λ must be as large as possible to maximize the signal obtained in the copper strips. A better approximation of the average induced charge can be obtained

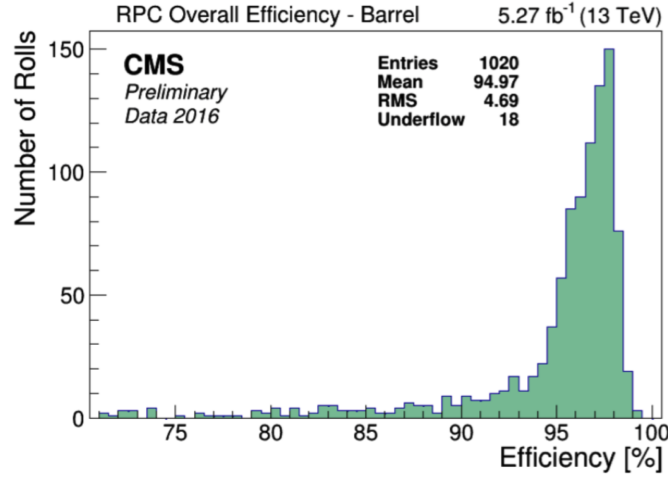


Figure 4.19: Efficiency (%) of each RPC in the CMS Barrel to provide reconstructed hits [188].

by using Monte Carlo simulations, where the fluctuations of the avalanche produced can also be taken into account.

RPC Performance

The RPCs are efficiently taking high-quality data. The efficiency is determined by taking the ratio of number of detected to the number of expected muon hits. This latest number is obtained from the extrapolation of muons reconstructed only in the muon system of the CMS experiment, without the RPC system. Segments (DT in the Barrel and CSC in the Endcap) that belong to a muon track, are selected and extrapolated to the plane of a given RPC. The detector unit is considered efficient if an RPC reconstructed hit is found within ± 2 strips from the position extrapolated from the DT/CSC segment [188]. Figures 4.19 and 4.20 show the efficiencies of the RPC chambers for both Barrel and Endcap, respectively.

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4.0.6 Trigger System and Data Acquisition

The amount of data that can be obtained by the LHC at each proton bunch crossing through the interaction point is about 1 MB. With a rate of crossings of 40 MHz, the amount of data can reach 40 TB every second. With the current storage technology is almost impossible to keep this amount of data [189]. The Trigger System reduces the number of events that are processed at each crossing of beams to about 100 events (100 MB) per second making a selection of the events that are best candidates for new physics [190]. The LHC has two trigger systems to accomplish this selection: The Level 1 Trigger (L1T) and the High Level Trigger (HLT)

L1 Trigger

The L1 Trigger is implemented in the detector electronics system to analyze the information that is available immediately after a proton-proton collision has occurred. The decision to accept or reject

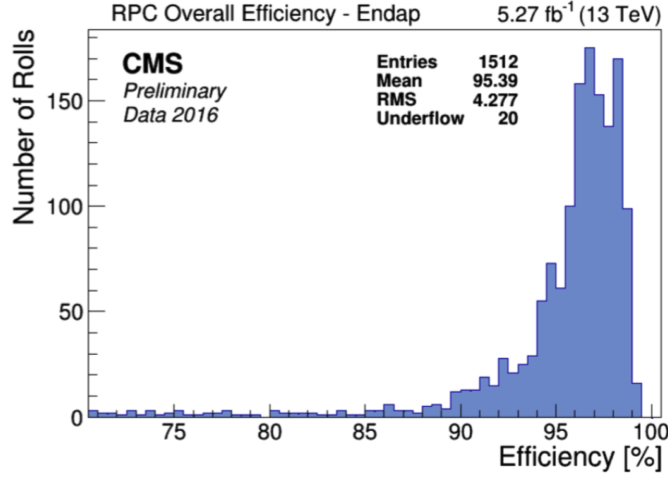


Figure 4.20: Efficiency (%) of each RPC in the CMS Endcap to provide reconstructed hits [188].

an event is based on data obtained by the electromagnetic calorimeter, the hadron calorimeter and the muon detectors. The data obtained by the tracker is not used because the algorithms used in this detector are too slow to serve as a trigger signal. The L1 Trigger reduces the rate of events to 50 KHz for low luminosity runs and 100 KHz for high luminosity runs. All data is stored temporarily while the trigger decision to accept or reject an event is taken. The time taken for the L1 Trigger to accept or reject the data coming from a particular proton bunch crossing is about $3.2 \mu s$. If the L1 Trigger accepts these events, they are processed by the next level of trigger (HLT). The L1 Trigger is divided into the calorimeter trigger (including the ECAL and HCAL) and the muon trigger.

The muon system plays a very important role for the L1 Trigger, because high transverse momentum muons provide a clean signal for many processes (including Supersymmetry). The L1 muon trigger is divided into three sections: DT, CSC and RPC. Each of these sections look for muons traces in the form of aligned impacts in each detector (as seen in Figure 4.1). Then the transverse momentum for the event is calculated and all the information is sent to the global muon trigger system. This system makes a correlation of the traces obtained in each muon detector and selects the best four muon candidates with the highest transverse momentum. Finally, this information is sent to the L1 Global Trigger.

The calorimeter trigger identifies from the region of the calorimeter the best four candidates of the following classes: electrons, photons, central jets, and τ -jets that are identified by the shape of the energy deposited in the calorimeter. Then, this information is passed to the L1 Global Trigger. Finally, the L1 Global Trigger makes the decision to accept or reject an event based on individual and collective analysis of all the events sent by the global muon trigger and the calorimeter trigger. This decision takes approximately $0.25 \mu s$ [189].

HLT Trigger

The High Level Trigger (HLT) is a software system implemented in a cluster of about one thousand processors. Its purpose is to select events that are candidates to the physics of interest of the experiment. This is accomplished by reconstruction and selection of events passing the L1 trigger reducing the event rate to 400 Hz. This trigger level performs a labeling of events. If they contain

one muon, two muons, jets, two jets, multiple jets, etc. This allows an additional labeling of the events according to the type of physics that may contain (Standard Model, SUSY, etc) [190]. There are three levels of analysis implemented in the HLT Trigger: The first level only takes into account the information coming from the muon system and the calorimeters. The second level adds the information from the tracker. Finally, in the third level the complete event information is available [189].

Chapter 5

Identification and Reconstruction Algorithms

Hadrons are complexed objects, made of quarks and gluons, each of them carrying out a fraction of momentum of the hadron, as described by the parton distribution functions. Therefore, the physics at hadron colliders is characterized by the presence of a large number of particles after the collision, each with its own characteristics. The knowledge of these features leads to the development of identification and reconstruction algorithms. These algorithms allow the identification of the particles within the data at software level. Once a particle is well identified by the algorithm it is called a “physical object”.

5.1 Particle Flow Algorithm

Most of the analyses in the CMS experiment use the particle flow (PF) event-reconstruction algorithm [191]. The goal of the algorithm is to combine the information from all subdetectors to identify and reconstruct each particle that is within the acceptance of the detector. These particles can come from different sources, such as hard interactions due to collisions in the center of the detector or cosmic rays. The algorithm performs the reconstruction and identification of all stable particles present in an event, classifying all of them into five mutually exclusive groups: muons, electrons, charged hadrons, photons and neutral hadrons. Then, this list of particles is used as an input in further algorithms, as if it came from an event generator. The list is used to make an estimate of E_T^{miss} , reconstruct jets, tau identification, etc. The traditional jet reconstruction methods use only the energy deposited in the calorimeters to perform the reconstruction, while the particle flow algorithm combines the information from all the subdetectors to reconstruct the particles individually, and then group them in jets. This procedure results in an improvement, not only for jet reconstruction, also in lepton identification and E_T^{miss} estimation (as shown in Figures 5.1 and 5.2). One of the main challenges for the algorithm is that most of the particles produced in a proton-proton collision have low transverse momentum. For example, the average transverse momentum carried by stable particles produced is of the order of 10 GeV/c, while a quark or gluon jet has a total energy of 500 GeV/c. This means that to distinguish between processes coming from the standard-model background and new physics it is necessary to identify and reconstruct accurately the largest possible number of particles, most of them with low transverse momentum. Another challenge for the algo-

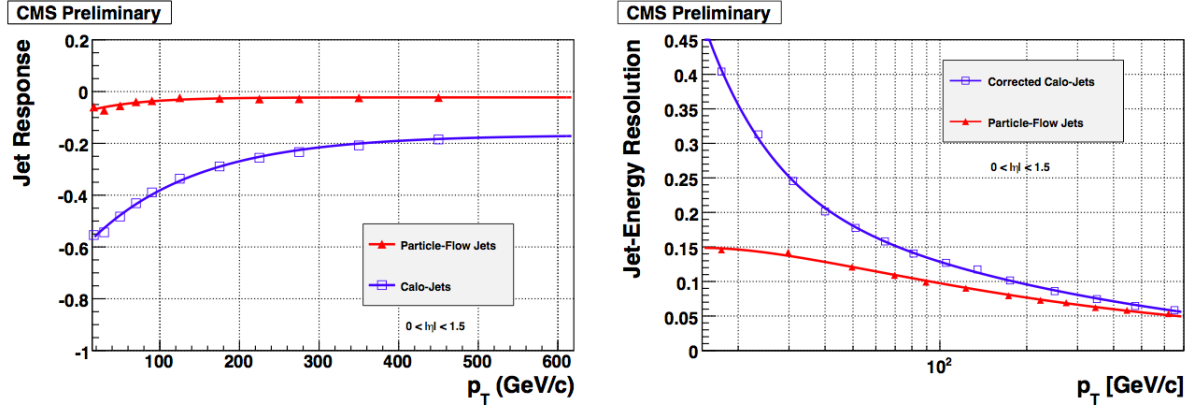


Figure 5.1: (*left*) Jet response for di-jet events in the barrel reconstructed with the particle flow (red triangles) and the calorimeters (blue open squares). (*right*) Jet resolution for di-jet events in the barrel reconstructed with the particle flow (red triangles) and the calorimeters (blue open squares) [191].

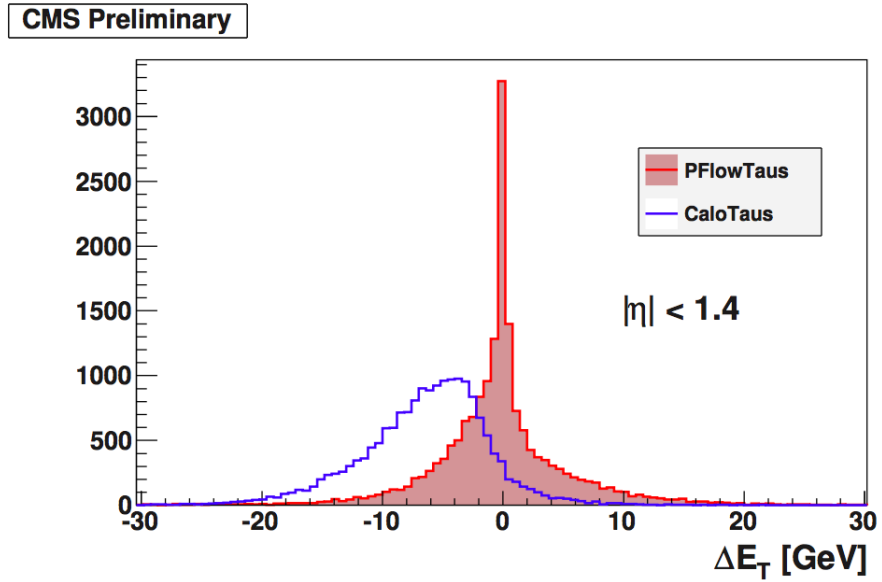


Figure 5.2: Energy resolution measured in the transverse plane for hadronic τ jets from Z decays, in the barrel, reconstructed using Particle Flow (filled), and using only calorimeter information (hollow) [191].

gorithm is to assure that the list of particles provided as an output has a good reconstruction efficiency and low number of fakes, even in events with a high density of particles, such as high pile up events. The particle flow algorithm performs the whole reconstruction process in three steps. First, it

reconstructs the PF elements, which are composed of tracks reconstructed in the tracker, the energy of clusters reconstructed based on the energy depositions in the calorimeters, and the reconstructed tracks in the muon system. In the second step, the algorithm performs an association of all the PF elements using the link algorithm (described in the following sections), creating linked elements. Finally, in the third step, all the individual particles are identified and reconstructed based on the content of the linked elements.

5.1.1 Iterative Tracking

The tracker is one of the main ingredients of the PF algorithm since its resolution, for particles with p_T up to 1 TeV scale, is much better than the resolution of the calorimeters. The tracker provides an accurate measurement of the direction of charged particles before any deviation produced by the magnetic field in their trajectory to the calorimeters. A high resolution and a low fake rate in the tracker are crucial for the PF algorithm, since most of the energy of the jets is recognized as coming from charged particles. To achieve the previous requirements, a very stringent criteria for the reconstruction of tracks is demanded in the first iteration, producing a very low fake rate, but at the same time the efficiency is reduced. In a second iteration, all the primary vertices are reconstructed and then all the hits used in the first iteration are removed and the reconstruction of the tracks is performed again with the remaining hits but now with a less restrictive criteria (as it is shown in Figure 5.3) and using only tracks that are compatible with primary vertices. The less stringent criteria produce an increase in the efficiency while the remotion of the hits of the first step maintain the fake rate low. Then, all the process is iterated again, also relaxing the constraints on the origin vertex. With this approach, charged particles with a transverse momentum as low as 150 MeV/c, and origin vertex greater than 50 cm from the beam axis, and with only three hits in the tracker are reconstructed with a fake rate of approximately 1% [191].

5.1.2 Particle Flow Clustering

There are two steps in PF algorithm. First, a seed selection is performed in the cells of the detector, in this case the calorimeters. The cell seed must have an energy greater than a given threshold, also the seed cell must have more energy than its neighbour cells as can be seen in Figure 5.4. It is important to note that the neighbours from a cell seed can not become a seed. Second, the topological clusters are created using as an input the seed cells from the previous step and adding cells with energy greater than a given energy threshold and that share one side in common with the cells already included in the cluster (see Figure 5.5). The energy thresholds are 80 MeV in the barrel and up to 300 MeV in the endcaps for the electromagnetic calorimeter (those values represent two standard deviations from the electronic noise). For the hadronic calorimeter the threshold is 800 MeV. The purpose of this algorithm is to have a high detection efficiency even for particles with low energy.

5.1.3 Link algorithm

In principle all the particles reconstructed by the PF algorithm should have traces in the different CMS subsystems, for example, hits in the trackers, clusters in the calorimeters and for the case of muons, muon hits in the muon chambers. This implies that all these elements need to be connected to make a full reconstruction of a signal particle trying to avoid also a potential double counting.

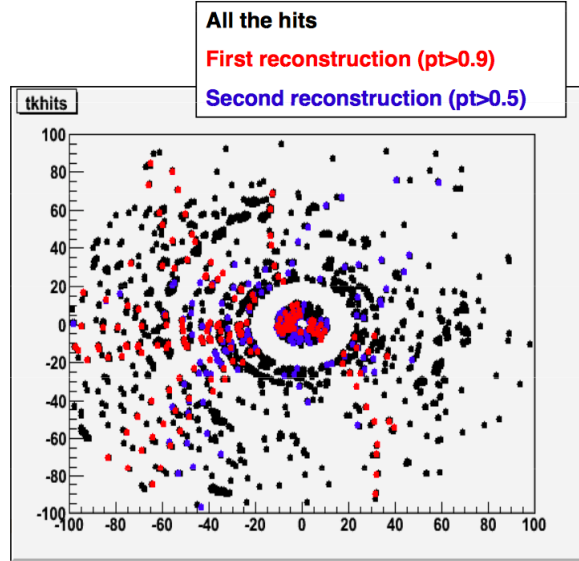


Figure 5.3: Iterative tracking algorithm, the hits with a stringent reconstruction criteria (in red) are removed, to perform again the reconstruction with a less stringent criteria (hits in blue). This procedure increases the reconstruction efficiency, while the fake rate is maintained low [192].

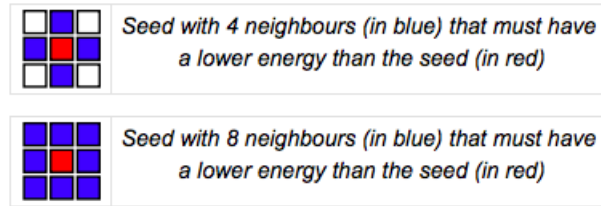


Figure 5.4: The selection of a seed for the particle flow clustering algorithm. The seed must have energy greater than its neighbours, and also energy greater than a given threshold [193]

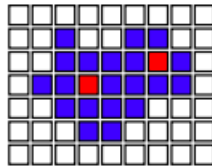


Figure 5.5: Example of a topological cluster created from cell seeds. A topological cluster can contain several seeds represented here in red [193]

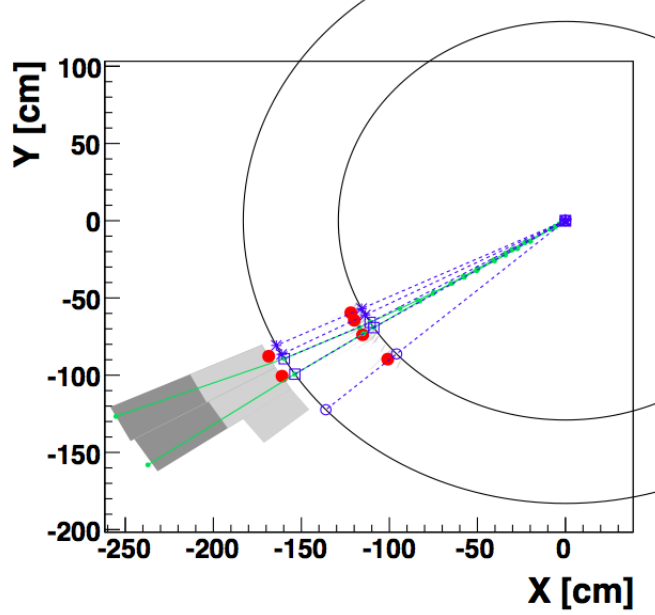


Figure 5.6: Event display of a simple hadronic jet in the X-Y plane, the Electromagnetic calorimeter and the Hadronic calorimeter are represented by the two circles centred around the interaction point (in blue) [191].

The connection decision is checked between two elements, and establish a distance between the two elements to give a quantification of the quality of the link. Then, the algorithm produces blocks that are elements related directly or indirectly between them. Due to the granularity of the experiment, those blocks commonly have one, two or three elements. This ensures that the performance of the algorithm has no dependency of the event complexity. An example of this is that a simple jet (as shown in Figure 5.6) has the same energy response and resolution than jets produced in QCD events, that normally contain a greater number of blocks.

5.2 JETS

Due to color symmetry in QCD, all observable particles have a net color charge of zero, therefore a quark or a gluon cannot be observed directly, they can only be observed in colorless final states. What the detector observes is the hadronization process of the quark or gluon, that produces a flow of hadrons in a localized area, called jet, which is detected by the tracker and the calorimeter system [171] (see Figure 5.7).

The goal of the reconstruction algorithms of jets in pp collisions is to be able to give an initial estimate of the energy and direction of the parton [169]. The idea of the algorithm is to group all particles originated from the initial parton into a single object, the jet [154]. The CMS experiment uses mainly the PF approach for the reconstruction of the jets, that means that all the particles of an event are identified and reconstructed individually using the information from all the subdetectors. This procedure is executed before performing any jet clustering algorithm. The hadronic

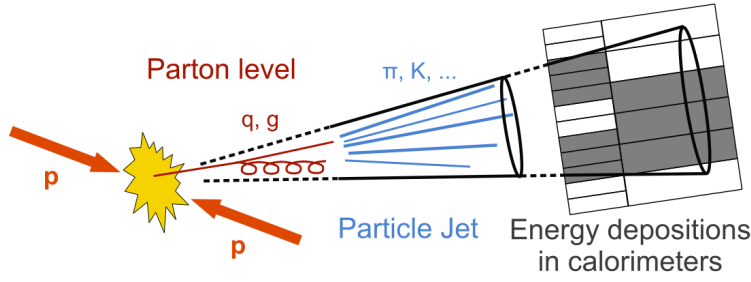


Figure 5.7: Sketch of the hadronization process after production of quarks and gluons in a proton-proton collision. [171].

and electromagnetic calorimeters are the main factors for the jet reconstruction performance. This performance depends on the interactions between hadrons and the calorimeter that is not linear. There are few algorithms in the CMS experiment that can be used for the reconstruction of jets:

- *Iterative Cone (IC)*: This is simpler and faster than the sequential recombination algorithms, commonly used for jet reconstruction in software-based trigger systems, within CMS, it is used in the HLT Trigger. Each cell of the calorimeter with energy higher than a threshold E_0 (in the case of CMS, $E_0 = 1\text{GeV}$) is considered as a seed to search for a cone. A jet is defined by adding the cells within the cone $R_{cone} = \sqrt{(\eta_i - \eta_{seed})^2 + (\phi_i - \phi_{seed})^2}$. The direction of the jet can be obtained with $\eta_{jet} = \sum_{i \in cone} E_{\perp i} \eta_i / E_{\perp jet}$, $\phi_{jet} = \sum_{i \in cone} E_{\perp i} \phi_i / E_{\perp jet}$, $E_{\perp jet} = \sum_{i \in cone} E_{\perp i}$. If the direction of the jet does not match the seed cell, the process continues, replacing the seed cell by the jet direction obtained. This process continues until a stable configuration for the jet is achieved [194].
- *Sequential recombination algorithms* : The first step is to define two distances: distances d_{ij} between particles i and j and d_{iB} between particle i and the beam (B). Objects with a minimum distance d_{ij} iteratively combined until only particles with distance $d_{ij} > d_{iB}$ remain. The resulting object is called a jet. These jets are then removed from the list of particles. Then, the distances are recalculated and the procedure is repeated until no particles are left. The difference between sequential recombination algorithms lies in the definition of the distance

$$d_{ij} = \min(k_{ti}^{2p}, k_{tj}^{2p}) \frac{\Delta_{ij}^2}{R^2}, \quad (5.1a)$$

$$d_{iB} = k_{ti}^{2p}, \quad (5.1b)$$

where $\Delta_{ij}^2 = (y_i - y_j)^2 + (\phi_i - \phi_j)^2$, k_{ti} is the transverse momentum, y_i is the rapidity, ϕ_i is the azimuth angle of particle i and R is the radius parameter. This radius parameter R serves to rescale the distance d_{ij} such that any pair of jets is separated by at least $\Delta_{ij} = R$. The value $p = 1$ corresponds to the kt algorithm [195, 196], $p = 0$ to the Cambridge/Aachen [197] and $p = -1$ to the anti-kt algorithm [198]. All three algorithms are infrared and collinear safe. This means that if one modifies an event by a collinear splitting or the addition of a soft emission, the set of hard jets that are found in the event should remain unchanged as can be seen in Figure 5.8. CMS and ATLAS use mainly the anti-kT algorithm although Cambridge-Aachen is used in several CMS analysis. The advantage for using anti-kT algorithm is due to the jet

area safety. The anti- k_T algorithm calculates first the distance for high- p_T particles and then includes low- p_T particles. This results in a slightly fluctuation in the jet area as it is shown in Figure 5.9.

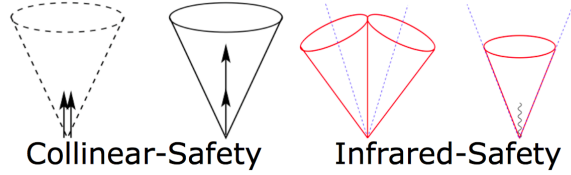


Figure 5.8: Sketch of the difference between collinear and infrared safe [171].

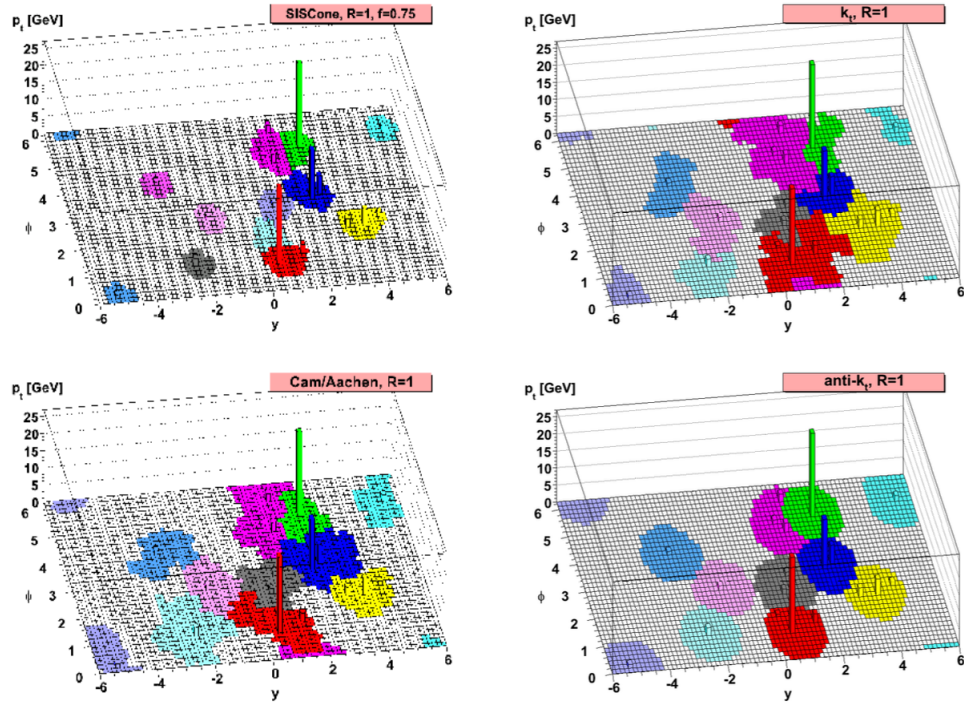


Figure 5.9: Jet reconstruction using the algorithms SIScone (top-left), k_T (top-right), Cambridge-Aachen (bottom-left) and anti- k_T (bottom-right). Jet reconstruction was performed with MC samples using the same R for all the algorithms. k_T and Cambridge-Aachen jet reconstruction show an irregular jet area, while SIScone reconstructs jets in smaller areas but identifies two different jets instead of the one grey jet [199].

The resulting jets with these algorithms are called raw jets because no correction has been applied to them. One of the corrections applied to the jets are due to the high number of pileup interactions, this is a source of complications during jet reconstruction, given the mixture of particles that are produced in different primary vertices. For this reason, it is necessary to include corrections to the

jets reconstructing algorithms. Typical jet energy resolutions at the central rapidities are 15-20% at 30 GeV, about 10% at 100 GeV, and 5% at 1 TeV [191].

5.2.1 Jet selection used in the analysis

Jets are required to have $p_T > 60$ GeV and $|\eta| < 5.0$. For the identification of jets the loose PF ID is used in this analysis. Table 5.1 shows the selection criteria used for the jet identification and reconstruction in this analysis. Figure 5.10 shows the η distribution of the highest p_T jet in a soft- μ + 1b + VBF dijet control sample enriched with $t\bar{t}$ background events, the good level of agreement observed provides further confidence in the reconstruction/identification of high- η jets.

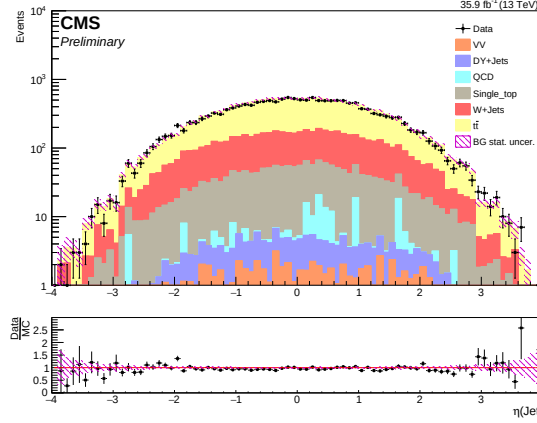


Figure 5.10: η distribution for jets in a $t\bar{t}$ control sample with one soft muon, 1 b-tagged jet, $E_T^{\text{miss}} > 250$ GeV, and $m_T(\mu, E_T^{\text{miss}}) < 110$ GeV.

Table 5.1: Required selections applied to the jets used in the analysis.

Selection	Cut
Jet p_T	≥ 60 GeV
$ \eta $	≤ 5.0
Jet Identification	Loose
ΔR	0.4

5.3 B tagging

Typically, the distinction between different jet flavors is challenging because the jet deposition signatures are similar. In particular, jets coming from up, down and strange quarks cannot be distinguished. On the other hand, b jets can be distinguished from other jets because bound states including b-quarks are long-lived ($\sim 10^{-12}s$) in comparison with bound states of other quarks. These bound b-quark states can travel a distance of the order of one millimeter before decaying. This distance allows the tracker to perform a secondary vertex reconstruction, since there is a relatively

good separation between the primary and the secondary vertex. One of the standard parameters used to discriminate between tracks coming from the second and primary vertex is called the impact parameter. The impact parameter of a track is defined as its distance to the primary vertex at the point of closest approach, as indicated in Figure 5.11. B-jets are constructed using as input objects

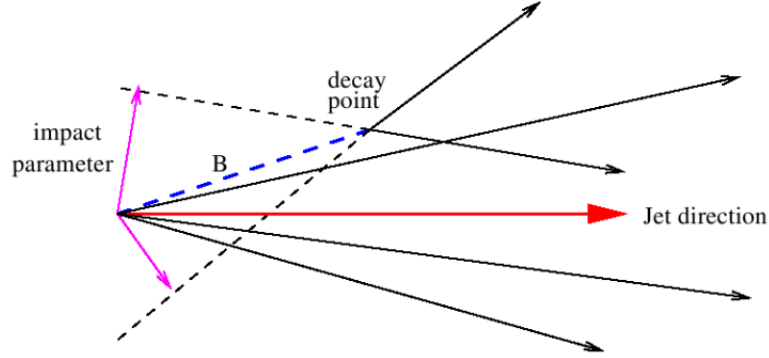


Figure 5.11: Illustration of the track impact parameter. The dashed blue line marks the flight distance of the B-meson. The extrapolations of the tracks that originate from its decay are indicated by the dashed black lines [200].

coming from the PF algorithm. These objects are matched together using the anti- k_T algorithm and a distance parameter of 0.5 to build the jets [201]. In principle, all the energies of the obtained jets need to be corrected to have a uniform response in η and an absolute calibration for the transverse momentum. Then, the kinematic properties of the objects inside the jets are used as an input for the different algorithms employed for the b-jet identification. The path of those objects is calculated in the tracking system of the CMS experiment with an iterative procedure using a standard Kalman filter method [202].

5.3.1 Combined Secondary Vertex algorithm

In our analysis, we used the Combined Secondary Vertex algorithm (CSV) for the identification of the b-jets. This algorithm has shown to have an excellent performance in the identification process [203]. The algorithm uses a multivariate analysis (MVA) using as input information: displaced tracks, secondary vertices, and the jet kinematics. The algorithm provides a discriminator for evaluating the compatibility of a given jet with a b-jet. This discriminator can have three different working points that are chosen according to the misidentification rate and the b-tagging efficiency.

- CSVL: loose working point, allows to select a high b-tagging efficiency (≈ 83 %) with a misidentification probability of 10 %.
- CSVM middle working point, allows to select a high b-tagging efficiency (≈ 69 %) with a misidentification probability of 1 %.
- CSVT tight working point, allows to select a high purity b-jet sample but with low b-tagging efficiency (≈ 49 %). The misidentification probability of is 0.1 %.

Table 5.2 shows the b-tagging efficiency for the three different working points of the CSVv2 discriminator.

The b-tagging algorithm was improved for run 2 (CSVv2) by including a multivariate algorithm that estimates the discriminator value using a neural network. This new algorithm allows an increase of 10 % in the b-tagging efficiency compared with the CSV algorithm used in run 1, as can be seen in Figure 5.12. The required selections applied to the b-jets used in our analysis are shown in Table 5.3.

Table 5.2: CSVv2 discriminator values and corresponding efficiency for the three working points. To determine the efficiencies, simulated events of $t\bar{t}$ produced in proton-proton collisions at $\sqrt{s} = 13$ TeV were used. Then b-jets with transverse momentum above 30 GeV were selected. Since the efficiency depends on the p_T and η of the jets, the numbers are presented just for illustration [204].

Working Point	CSVv2 discriminator	b-tagging efficiency	Misidentification probability
Loose	≥ 0.5426	≈ 83 %	10 %
Medium	≥ 0.8484	≈ 69 %	1 %
Tight	≥ 0.9535	≈ 49 %	0.1 %

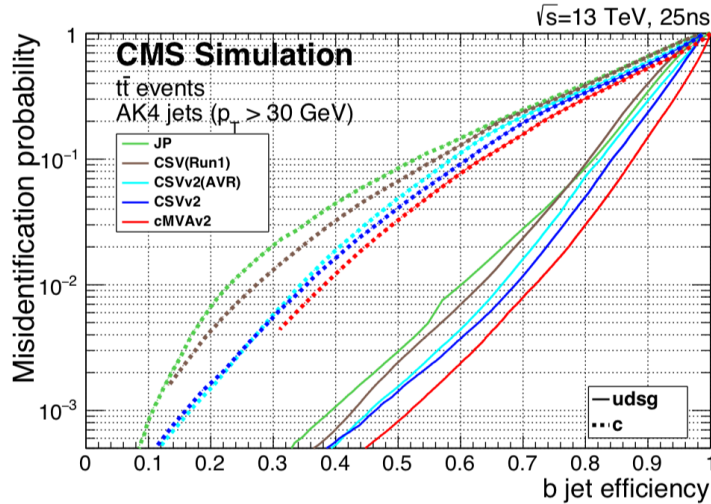


Figure 5.12: Performance of the b-jet identification efficiency algorithms demonstrating the probability for non-b jets to be misidentified as b jet as a function of the efficiency to correctly identify b jets. To determine the efficiencies, simulated $t\bar{t}$ events produced in proton-proton collisions at $\sqrt{s} = 13$ TeV were used. Although cMVA2 algorithm shows the best performance, the CSVv2 algorithm is used for many analysis and shows an improvement compared with the CSV algorithm used in Run I. The performance in this figure serves as an illustration since the b-jet identification efficiency depends on the p_T and η distribution of the jets in the topology. [204].

Table 5.3: Required selections applied to the b-jets used in our SUSY VBF analysis.

Selection	Cut
Jet p_T	≥ 30 GeV
$ \eta $	≤ 2.4
B-tagging identification algorithm	CSVv2
Working Point	Medium (0.8484)

5.4 MET

The CMS detector has a solid angle coverage close to 4π , although it is not fully closed because it is impossible to cover the openings where the proton beams come. The calorimeter system covers a range of $|\eta| < 5$ ($0.77^\circ < \theta < 179.23^\circ$). Therefore the total energy balance can not be used, since low momentum particles escape detection in the beam direction. Nevertheless, these particles have small transverse momentum, this fact allows to perform a rather precise measurement of the momentum conservation in the plane perpendicular to the direction of the beams. As a result, any significant imbalance in the transverse momentum measured in the calorimeter is an indicative of the production of a particle in the collision that interacts weakly. Within the SM, an imbalance of this kind in the transverse momentum indicates the presence of a muon or a neutrino. But the momentum of the muon can be accurately measured by combining the central tracker and the muon system. And thus the transverse momentum can be corrected including the momentum of the muon. The unique SM particle that escapes detection is the neutrino. In SUSY models that conserve R parity, the LSP is stable, massive, only interacts weakly and additionally is always produced in decays of heavier SUSY particles. The presence of LSP particles in SUSY events contribute to the missing transverse energy, which is therefore an important observable to identify SUSY events [154]. The missing transverse energy is determined by the vector sum of the transverse energy deposited in the cells of the calorimeter:

$$\vec{E}_T^{miss} = - \sum (E_n \sin \theta_n \cos \phi_n \hat{i} + E_n \sin \theta_n \sin \phi_n \hat{j}) \quad (5.2)$$

$$\vec{E}_T^{miss} = E_x^{miss} \hat{i} + E_y^{miss} \hat{j} \quad (5.3)$$

The missing transverse energy is extremely sensitive to detector malfunctioning and particles that reach regions that have deficient detection. Any failure or malfunctioning on the detector results in an artificial imbalance in transverse energy, which would be a false signal of a SUSY event. The use of this variable in the present work is to separate SM events from SUSY events. In our analysis we select events with $E_T^{miss} > 250$ GeV. This requirement is motivated in Section 6.4.2.

5.5 Muons

Muon reconstruction is performed in three steps: Local reconstruction (muon chambers are used), standalone reconstruction (the muon system is used) and global reconstruction (the tracker and the muon system are used). The local reconstruction uses the position of the tracks left on the muon

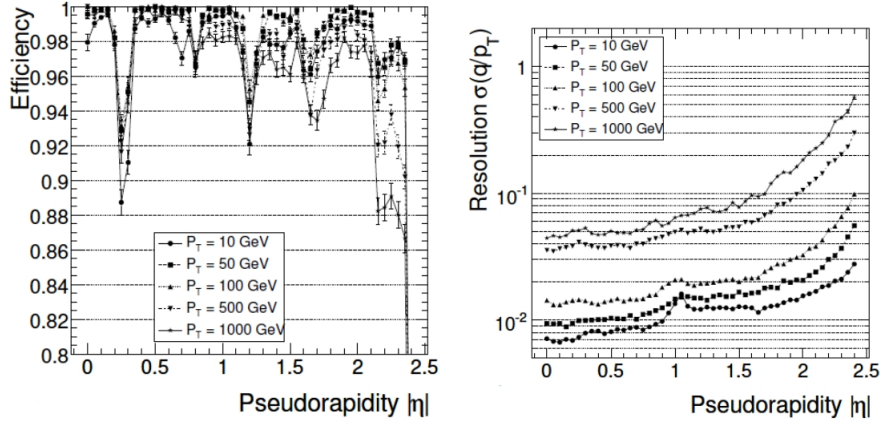


Figure 5.13: Reconstruction efficiency (left) and momentum resolution (right) as a function of pseudorapidity. [154].

chambers and forms what is called segments. These segments are used to generate state vectors (called track segments). This state vectors consist of position, direction and a rough estimate of the muon transverse momentum. The CSCs and DTs are used to determine position, the RPCs are also used but they have lower spatial resolution. The RPCs are used to complement CSCs and DTs in regions containing geometric coverage problems, in particular, regions where there is overlapping between barrel and endcap detectors. The state vectors obtained from the local reconstruction are used as the input of the standalone reconstruction. The muon trajectories are constructed starting with the most inner chambers, and then including outer chambers using the Kalman filter technique [154]. To reject tracks in the detectors coming from cascades, delta rays and pair production a cut is applied in the χ^2 variable that is obtained from a fit to the muon trajectory reconstruction. Once this process is completed, a stand Alone Muon is obtained [62].

The global reconstruction includes the information coming from the inner tracker system to extend the reconstructed muon trajectory. In this process the input information is the standalone reconstruction of the muon; an extrapolation of the muon trajectory is performed starting from the innermost muon station and ending with the outermost region of the inner tracker system. The region of interest in the tracker is determined by the extrapolation of the muon track, taking into account the energy lost by the muon in the material, effects of multiple scattering and uncertainties in the extrapolation. Within the regions of interest in the inner tracker system, regional seeds are created using two tracks from different layers. Then, a track reconstruction algorithm runs transforming each seed in a set of trajectories from the inner to the outer. Ambiguities in the trajectories coming from the same seed are solved based on the multiplicity of tracks and several tests on the χ^2 value. In the last step, a fit is performed on the reconstructed tracks using the hits that come from the standalone muon reconstruction. The final muon candidates are selected based on a cut in χ^2 . The reconstruction efficiency and momentum resolution for the global muons is shown in Figure 5.13. The reconstruction efficiency is typically 95-99%, except for pseudorapidity regions that lie just between two DTs ($\eta = 0.25$ and $\eta = 0.8$) or in the transition region between DTs and CSCs ($\eta = 1.2$). For low momentum muons, the muon resolution is about 2% and it is mostly obtained by the resolution in the silicon tracker.

5.5.1 Muon reconstruction and Identification improvements

A small reduction of muon reconstruction efficiency in the tracker, which worsened with pileup, was observed in 2012 data. In order to recover the efficiency, two additional muon- specific tracking iterations have been implemented for Run2:

1. An Outside-in iteration, seeded from the muon system, designed to recover the missing muon-track in the tracker.
2. An Inside-Out iteration designed to re-reconstruct muon-tagged tracks with looser requirements to improve the hit-collection efficiency.

In addition to recovering the small inefficiency, the new algorithm has also improved the muon track quality, in particular the number of hits per track. This is reflected in the identification efficiency of standard selections as the Tight Muon selection [205]. Tag-and-Probe efficiencies for Tight Muon ID on $Z \rightarrow \mu\mu$ events from a subsample of data recorded in the last part of the 2012 Run, reconstructed with both the Run1 and Run2 reconstruction algorithms are shown in Figure 5.14. The Run2 reconstruction improves the average efficiency by 1-2%, in particular in events with high multiplicity of proton- proton interactions per bunch crossing (pileup) [205].

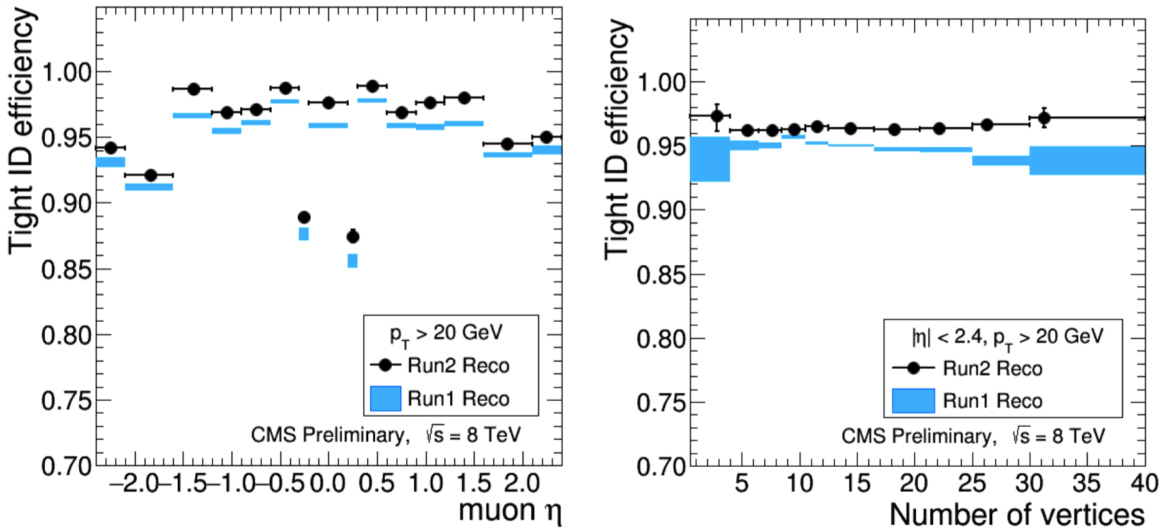


Figure 5.14: Efficiency as a function of muon pseudorapidity (left) and number of reconstructed vertices (right) for muon $p_T > 20$ GeV and $|\eta| < 2.4$ [205].

5.5.2 Muon trigger improvements in HLT

The muon triggers were updated with respect to the 2012 configuration. Single muon and Double muon trigger paths were improved, both in the online reconstruction and in the selection and quality cuts applied [181]. In particular, for the Single muon triggers, the tracking algorithms and pileup mitigation techniques were improved. Isolation computation was updated from single isolation cut (combining tracker and calorimeter information) to sequential isolation cuts, tuned independently for tracker and calorimeter components. For the Double muon triggers, a loose track based relative isolation cut was also added to both legs of the high p_T trigger, to reduce rates and keep thresholds identical to the ones used at the end of Run1.

5.5.3 Dimuon invariant mass spectrum

Figure 5.15 shows the dimuon invariant mass spectrum in the vicinity of the J/ψ meson and $\Upsilon(\text{ns})$ meson respectively with 2016 data. These plots show the high accuracy for the reconstruction of the different dimuon invariant mass resonances using the muon system of the CMS experiment.

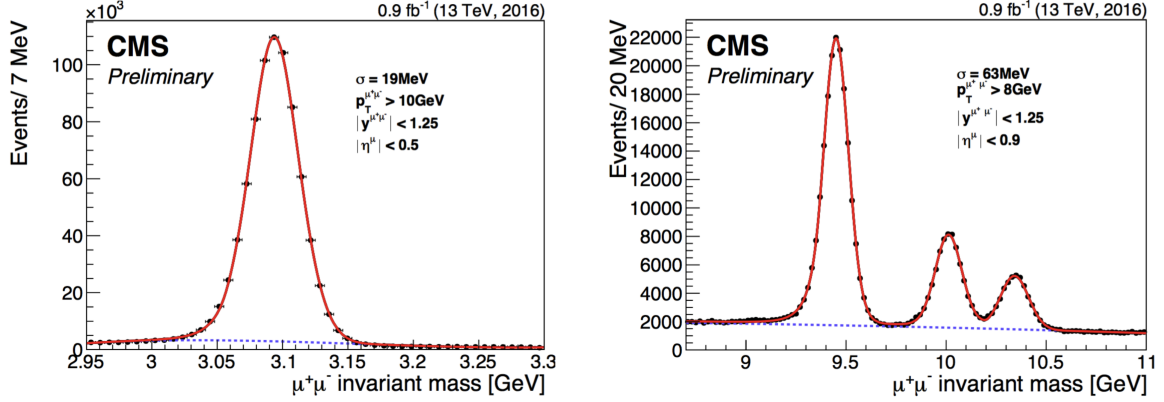


Figure 5.15: Dimuon invariant mass spectrum in the vicinity of the J/ψ meson (left) and $\Upsilon(\text{ns})$ meson (right) [206].

5.5.4 Efficiency estimation

Efficiencies are computed by means of the tag-and-probe method exploiting $Z \rightarrow \mu^+\mu^-$ resonances. The method requires a muon pair consisting of a tagged muon that triggers the event, with a tight requirement, and a probe muon with a loose selection used for the efficiency measurement [206]. Figures 5.16 shows Loose ID efficiency and Tight ID efficiency as a function of $|\eta|$. The efficiency is close to one for the two identification selections for $|\eta| < 2.4$.

5.5.5 Muon selection used in the present analysis

Muons are required to have a p_T requirement of [8-40] GeV and $|\eta| < 2.1$. An additional background reduction is obtained through a relative isolation requirement defined as:

$$\text{Relative Isolation} = \frac{\sum p_T^{\text{charged hadrons}} + \sum p_T^{\text{neutral hadrons}} + \sum p_T^{\gamma} - 0.5 \times \sum p_T^{\text{charged hadrons from PU}}}{p_T^{\mu}} \quad (5.4)$$

where the sums run over the charged hadron candidates, neutral hadrons and photons, within a ΔR cone around the muon direction. In this analysis, the isolation requirement is $\Delta R < 0.4$. We use the "Tight" identification working point and relative isolation < 0.25 . Table 5.4 shows the selection criteria used for the μ identification.

5.6 Electrons

The reconstruction of electrons in the CMS experiment requires to relate the energy deposited on a cluster of the electromagnetic calorimeter to a track from the tracker. The electromagnetic

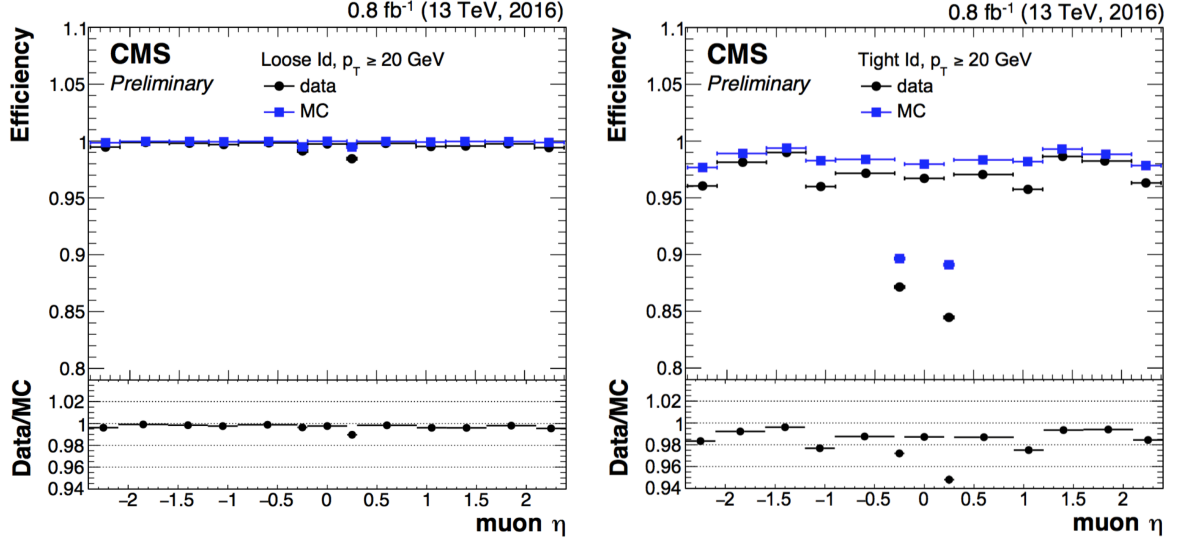


Figure 5.16: Loose ID efficiency vs $|\eta|$ (left) and Tight ID efficiency vs $|\eta|$ (right) computed with the tag- and-probe method using collision data at 13 TeV and 25 ns bunch spacing [207].

Table 5.4: Required selections applied to the muons used in the present analysis.

Selection	Cut
Muon p_T	[8-40] GeV
$ \eta $	≤ 2.1
Muon Identification	Tight
Relative Isolation	0.25
ΔR used for Isolation	< 0.4

calorimeter has been designed to collect all the energy coming from an incident electron. A cluster is defined as a region centered on a high energy deposited crystal that extends from that point in the directions η and ϕ until the energy falls below certain range, or until the energy grows back. The energy of the electron is then taken from the energy of the cluster. When for some reason there is an indication that the energy of the cluster has been measured wrongly, the electron energy is taken from the measured momentum of the track. The track is used to derive the electron charge and to reject the background by applying quality cuts in the momentum, the energy of the cluster, and also making matches between the cluster energy and the track in the tracker.

The CMS reconstruction code provides three different levels in the identification of electrons: robust, loose and tight. The robust identification is designed to be simple and insensitive to the expected uncertainties in the alignment and calibration of the detector. Loose and tight identification involve specific cuts that take into account the energy of the cluster, the momentum at the entrance, and the classification of the electron. The electrons that are identified as tight are those which pass the loose identification and are also identified as robust. Table 5.5 shows the selection criteria used for the electron identification in the present analysis. Since we are focused on final states with exactly

one well-identified muon, events that contain an electron that satisfies the selections displayed in Table 5.5 are vetoed.

Table 5.5: Required selections applied to the electrons used in the present analysis.

Selection	Cut
Electron p_T	>10 GeV
$ \eta $	≤ 2.1
Electron Identification	Medium
Relative Isolation	0.25
ΔR used for Isolation	<0.4

5.7 Taus

The tau leptonic decays cannot be reconstructed because it is impossible to distinguish the leptons produced in the tau decays from leptons originating from hard interactions in proton-proton collisions. Figure 5.17 shows the Feynman diagrams for the hadronic and leptonic decay of the tau. Then, only tau hadronic decays can be identified directly using the hadronic calorimeter. The hadronic tau reconstruction also represents a challenge since the signal in the calorimeter is very similar to that produced by QCD jets, with the difference that the latter has a cross section eight orders of magnitude higher. An additional challenge comes from the fact that taus have neutrinos in their decay modes, then the tau kinematic variables cannot be completely reconstructed since neutrinos escape detection.

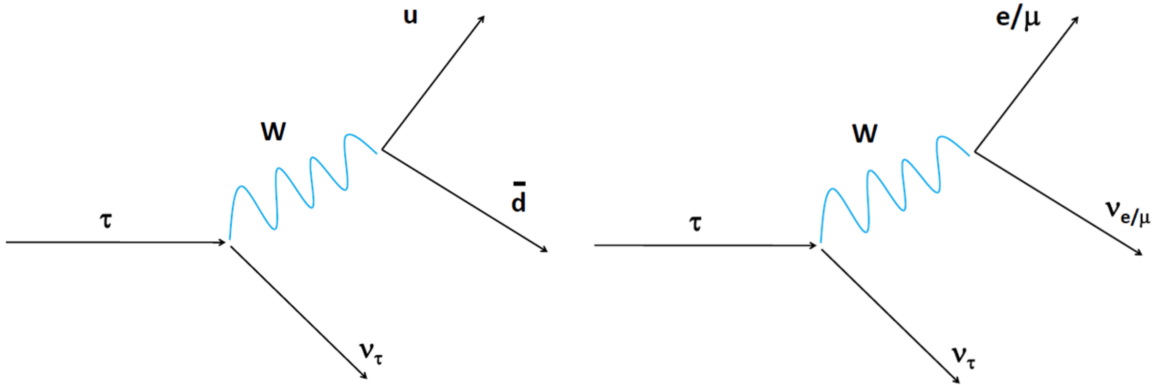


Figure 5.17: Feynman diagrams for tau decays. Hadronic decay with a branching fraction of 64.8% (left) and leptonic decay with a branching fraction of 35.2% (right) [23].

5.7.1 HPS algorithm

In most analyzes of the CMS experiment, the Hadron-Plus-Strips (HPS) algorithm [208] is used for the reconstruction and identification of the taus. The algorithm is executed in two steps. First, a

search of neutral and charged PF objects, that are compatible with the decay modes of the tau, is executed. In the next step, several discriminators are applied to the reconstructed taus to distinguish them from processes that resemble a tau, mainly, QCD jets, and electrons and muons with a lower contribution. All this with the ultimate goal of lowering the misidentification rate.

Table 5.6 shows the main decay modes and branching fraction for the tau. According to the table, the main hadronic tau decays are composed of one neutrino, neutral pions, and charged hadrons that are mainly pions, and kaons, in a much smaller branching fraction. This algorithm uses jets reconstructed by the PF algorithm as input parameters. With this algorithm, the different tau decay modes are reconstructed individually. This is achieved by discriminating between the neutral and charged components that make up the incoming jets. The PF algorithm facilitates this task since it reconstructs and identifies all the particles that make up the jets individually, to then use different clustering algorithms to form the jets.

Table 5.6: Main decay modes and branching fractions for negative taus [209].

Final State	Branching Fraction [%]
$e^- \bar{\nu}_e \nu_\tau$	17.8
$\mu^- \bar{\nu}_\mu \nu_\tau$	17.4
$h^- \nu_\tau$	11.5
$h^- \pi^0 \nu_\tau$	25.9
$h^- \pi^0 \pi^0 \nu_\tau$	9.5
$h^- h^+ h^- \nu_\tau$	9.8
$h^- h^+ h^- \pi^0 \nu_\tau$	4.8
Others	3.3

To search for charged hadrons, the HPS algorithm uses the information from the Tracker and the Calorimeter System. Additionally, energy depositions that are consistent with traces left by neutral pions are sought. There is a high probability that a photon that comes from a pion decay ($\pi^0 \rightarrow \gamma\gamma$) becomes an electron-positron pair due to the photon interaction with tracker material. Then due to the strong magnetic field produced within the solenoid, the electron and the positron are separated in opposite directions, resulting in elongated energy depositions in the ϕ direction (see Figure 5.18). Then, the HPS algorithm looks for ECAL strips in the $\eta - \phi$ plane that are consistent with a trace left by a π^0 . Finally, these strips are combined with the charged hadrons reconstructed to build possible tau decay modes.

5.7.2 Selection criteria applied for tau identification

Table 5.7 shows the selection criteria used for the tau identification in the present analysis. Since we are focused on final states with exactly one well-identified muon, events that contain a tau that satisfies the selections displayed in Table 5.7 are vetoed.

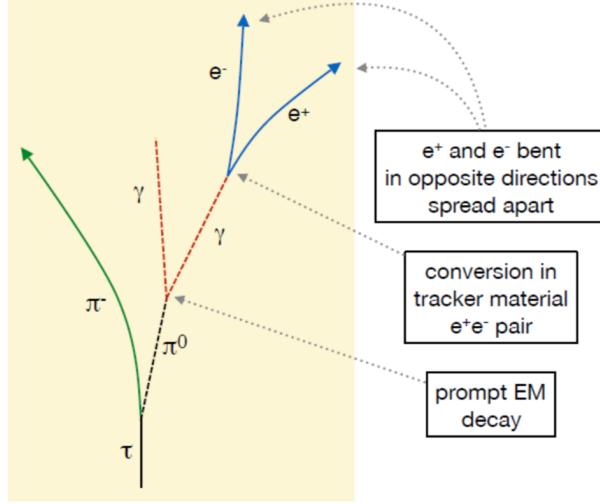


Figure 5.18: Schematic view of the tau decay to $\pi^-\pi^0\nu_\tau$. This decay represents $\sim 57.5\%$ of all hadronic tau decays.

Table 5.7: Required selections applied to the taus used in the present analysis.

Selection	Cut
Tau p_T	>20 GeV
$ \eta $	≤ 2.1
Discrimination prong type	1 HPS

5.8 Summary of physical variables

5.8.1 Transverse mass, m_T

The use of the m_T variable is motivated by the kinematics of a decay in two bodies as shown in figure 1. If we consider that we have a massive particle of mass m that decays to two particles 1 and 2. The transverse mass would be defined by the equation:

$$m_T^2 = m^2 + p_x^2 + p_y^2 = E^2 - p_z^2 \quad (5.5)$$

Where the z-axis is in the direction of the beam and p_x and p_y are the momentum projection on the plane defined by the x-axis and the y-axis, perpendicular to the direction of the proton beam. The transverse energy for this particle is defined as:

$$E_T = E \frac{|\vec{p}_T|}{|\vec{p}|} = \frac{E}{\sqrt{E^2 - m^2}} |\vec{p}_T| \quad (5.6)$$

where the transverse momentum is defined as

$$\vec{p}_T = (p_x, p_y) \quad (5.7)$$

By using the above definitions (Equations 5.5 and 5.6) and the four-momentum of particles 1 and 2 (p_1, p_2 , respectively), we can write the mass for the two-particle system as:

$$M_{12}^2 = (p_1 + p_2)^2 = p_1^2 + p_2^2 + 2p_1 p_2 = m_1^2 + m_2^2 + 2(E_1 E_2 - \vec{p}_1 \cdot \vec{p}_2) \quad (5.8)$$

$$M_{12}^2 = m_1^2 + m_2^2 + 2 \left(E_{T,1} \frac{\sqrt{p_{1,x}^2 + p_{1,y}^2 + p_{1,z}^2}}{p_{T,1}} E_{T,2} \frac{\sqrt{p_{2,x}^2 + p_{2,y}^2 + p_{2,z}^2}}{p_{T,2}} - \vec{p}_{T,1} \cdot \vec{p}_{T,2} - p_{z,1} p_{z,2} \right) \quad (5.9)$$

$$M_{12}^2 = m_1^2 + m_2^2 + 2 \left(E_{T,1} E_{T,2} \sqrt{1 + p_{1,z}^2 / p_{T,1}^2} \sqrt{1 + p_{2,z}^2 / p_{T,2}^2} - \vec{p}_{T,1} \cdot \vec{p}_{T,2} - p_{z,1} p_{z,2} \right) \quad (5.10)$$

now taking the transverse projection of the system (with $p_{1,z} = p_{2,z} = 0$) we obtain:

$$(M_{12}^2)_T = m_1^2 + m_2^2 + 2(E_{T,1} E_{T,2} - \vec{p}_{T,1} \cdot \vec{p}_{T,2}) \quad (5.11)$$

Note that an identical result can be obtained if we define the transverse mass as the one obtained from the relativistic equation of energy, but just taking into account energy and momentum of the daughter particles in the transverse plane, as previously defined:

$$M_T^2 = (E_{T,1} + E_{T,2})^2 - (\vec{p}_{T,1} + \vec{p}_{T,2})^2 \quad (5.12)$$

The previous expression is defined as the transverse mass. For the case we study, we have one lepton and a neutrino from the decay of a W boson, or a lepton and a neutralino from an slepton decay. The minimum p_T we are considering for the lepton is 8 GeV, still much greater than the mass of any SM lepton (the tau is the one with the highest mass, which is about 1.8 GeV, still a factor of 4 less than the minimum momentum considered). So, a good approximation is $|\vec{p}_T| \approx E_T$, in other words

we neglect the masses of the daughter particles, because they are small. With these approximations we obtain:

$$M_T^2 = 2E_{T,1}E_{T,2}(1 - \cos \phi) \quad (5.13)$$

Where ϕ is the angle between the two daughter particles in the transverse plane. Note that the distribution of M_T has a maximum $M_T^{\max} = M$ that is the invariant mass of the system. In this way, we can define the transverse mass m_T for the lepton- E_T^{miss} system as:

$$m_T = \sqrt{2 \cdot p_T^\ell \cdot E_T^{\text{miss}} \cdot (1 - \cos \Delta\phi_{\ell, E_T^{\text{miss}}})}, \quad (5.14)$$

where $\Delta\phi_{\ell, E_T^{\text{miss}}}$ is the azimuthal angle between the charged lepton's transverse momentum p_T^ℓ and the E_T^{miss} direction [210].

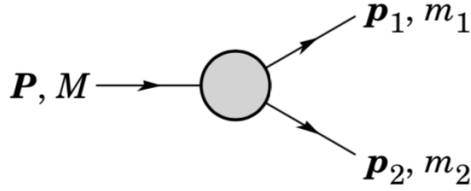


Figure 5.19: Variable definitions for a two body decay [13].

5.8.2 Dijet mass

The mass of any particle follows the relativistic equation of energy:

$$m^2 = E^2 - |\vec{p}|^2. \quad (5.15)$$

In the case of two jets, we add their energy and momentum and obtain a mass for the dijet system (invariant mass):

$$m_{jj}^2 = (E_{j_1} + E_{j_2})^2 - |\vec{p}_{j_1} + \vec{p}_{j_2}|^2 \quad (5.16)$$

Performing the algebra from the previous equation we arrive at:

$$m_{jj}^2 = m_{j_1}^2 + m_{j_2}^2 + 2(E_{j_1}E_{j_2} - \vec{p}_{j_1} \cdot \vec{p}_{j_2}) \quad (5.17)$$

Since we are considering relativistic jets ($p_{T_{\text{jet}}} > 60 \text{ GeV}$), $E_j \gg m_j$ and therefore

$$E_j = |\vec{p}_j| \quad (5.18)$$

The invariant mass of the dijet system simplifies to:

$$m_{jj}^2 = 2p_{j_1}p_{j_2}(1 - \cos(\theta_2 - \theta_1)) \quad (5.19)$$

In high energy physics we do not use the polar angle θ , it is better to use pseudorapidity (η), given that the difference in pseudorapidity is Lorentz invariant. After performing the corresponding algebra, the dijet invariant mass can be written as:

$$m_{jj} = \sqrt{2 \cdot p^{j_1} p^{j_2} \cosh(\Delta\eta(j_1, j_2))} \quad (5.20)$$

The use of this variable is motivated by the eta distribution for the jets that come from the signal in comparison with the jets in the background, as shown in Figure 5.20. The jets in SM background events are mostly central and have small $\Delta\eta(j_1, j_2)$, while signal events are characterized by more forward-like jets with large $\Delta\eta(j_1, j_2)$. Additionally, the outgoing partons in VBF processes must carry relatively large p_T since they must have enough momentum to produce the SM vector bosons as well as keeping their forward trajectories, implying large dijet invariant masses, which motivates the application of a high m_{jj} requirement to achieve a separation between signal and background events.

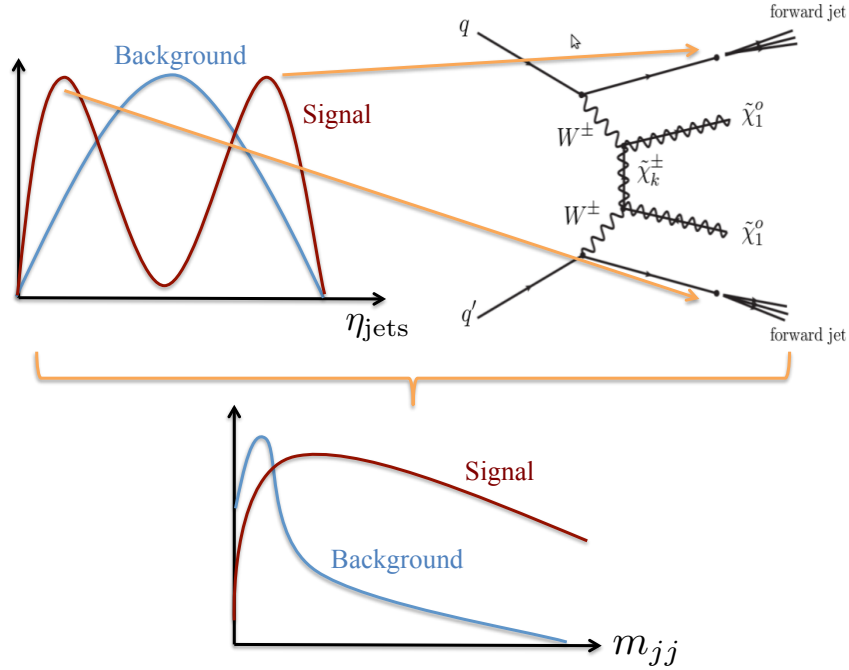


Figure 5.20: m_{jj} is motivated by the high $\Delta\eta$ and p_T of the VBF jets.

5.8.3 PFMET, H_T and MHT

PFMET makes use of the PF algorithms to reconstruct the momenta of individual particles.

$$\sum_i \vec{p}_T^{\text{detected}} + \sum \vec{p}_T^{\text{no-detected}} = 0$$

$$\sum \vec{p}_T^{\text{no-detected}} = - \sum_i \vec{p}_T^{\text{detected}}$$

$$\Rightarrow \vec{p}_T^{\text{miss}} = - \sum_i \vec{p}_T^{\text{detected}} \quad (5.21)$$

where the index i runs over all particle flow candidates. E_T^{miss} is the magnitude of $\vec{p}_T^{\text{miss}}$. The deposits coming from the Hadronic forward calorimeters (HF) are included in the calculation of E_T^{miss} and type-1 corrections are applied [211]. H_T is defined addition of the magnitudes of all jet momenta: $H_T = \sum \vec{p}_T$ where the sum is carried out over jets with $p_T > 50\text{GeV}$ and $|\eta| < 2.5$ [212]. MHT means “Missing H_T ” ($\vec{H}_T$), is defined as the vector sum: $\vec{H}_T = - \sum \vec{p}_T$ where the sum is over jets with $p_T > 30\text{GeV}$ and $|\eta| < 5$ [212]. The definition of these variables is similar to E_T^{miss} (see Figure 5.21). E_T^{miss} , H_T and MHT are expected to have higher values in events with SUSY particles, in comparison with SM particles. In the case of H_T , jets are taken into account within the acceptance of the tracker ($|\eta| < 2.5$) while in the case of MHT, the reconstructed jets are of less quality because a larger η region is included in its definition ($|\eta| < 5.0$).

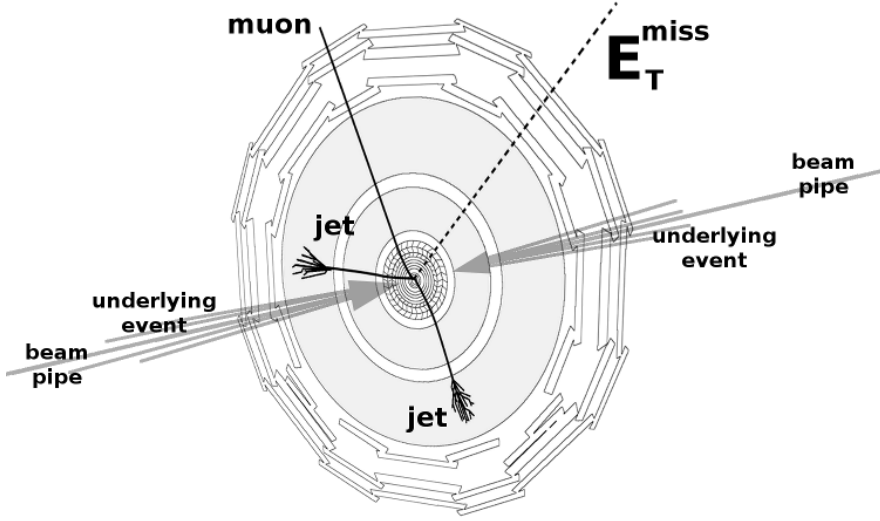


Figure 5.21: Scheme of an event with reconstructed E_T^{miss} [213].

Chapter 6

Analisis

As described in Section 2.9, the VBF SUSY topology has far-reaching sensitivity for electroweak SUSY searches in the compressed spectra parameter space than many other analysis techniques available at the moment. An experimental confirmation of this fact is the VBF SUSY search for electroweakinos performed by the CMS experiment at $\sqrt{s} = 8$ TeV using a proton-proton collision data sample with an integrated luminosity of 19.7 fb^{-1} [20]. Exclusion limits (at 95% CL) for the pair production cross sections of charginos and neutralinos with two associated jets were established in this study by interpreting the non-event-excess results within the MSSM. An interesting case of consideration is the $\tilde{\tau}$ dominated scenario (where the third generation of sleptons are the lightest among the three families) and assuming $m(\tilde{\tau}) < m(\tilde{\chi}_1^\pm)$, an upper limit of 170 GeV for the masses of $\tilde{\chi}_2^0/\tilde{\chi}_1^\pm$ were established for the compressed case where the mass difference of $\tilde{\chi}_2^0/\tilde{\chi}_1^\pm$ and LSP is 50 GeV. This mass gap of 50 GeV with light $\tilde{\tau}$ was not previously excluded by any other CMS analysis technique and proved the potential of the VBF tagging to increase the sensitivity of SUSY electroweak searches. The 8 TeV VBF SUSY analysis exploded the fact that for light sleptons and mass gaps between $\tilde{\chi}_1^\pm/\tilde{\chi}_2^0$ and the LSP, the $\tilde{\chi}_1^\pm$ ($\tilde{\chi}_2^0$) would have preferentially a decay chain as $\tilde{\chi}_1^\pm \rightarrow \tilde{\ell}^\pm \nu_\ell \rightarrow \ell^\pm \nu_\ell \tilde{\chi}_1^0$ ($\tilde{\chi}_2^0 \rightarrow \tilde{\ell}^\pm \ell^\mp \rightarrow \ell^\pm \ell^\mp \tilde{\chi}_1^0$) as indicated in Figure 6.1. The compressed mass scenario is characterized by the SUSY parameter space where the mass of the mother particle is very close to the sum of the masses of their daughter particles such that the daughter particles are produced with very low kinetic energy. Therefore, the leptons in the final state have a very soft p_T spectrum and consequently very difficult to identify. Even though, the diagrams of Figure 6.1 show final states with two or three leptons, it is natural to expect reconstructed events with less number of leptons due to the soft-lepton reconstruction inefficiency. In the VBF SUSY 8 TeV analysis two leptons in the final state were studied with final state topology as: VBF jj + ℓ + ℓ + MET. An interesting analysis, that could potentially increase the physic sensitivity, is the VBF SUSY with one lepton in the final state (VBF jj + ℓ + MET), where ℓ could be either electron, muon or tau with some extra advantages that we can profit such as a higher center of mass energy ($\sqrt{s}=13$ TeV) that reflects as higher cross sections, and higher luminosity (the CMS 2016 proton-proton collision data sample corresponds to a luminosity of 35.86 fb^{-1}). The goal of the present thesis document is to describe all the analysis details for the event topology VBF jj + μ + MET using the 2016 proton-proton collision data sample recorded by the CMS experiment. The analysis of electron and tau channels were performed by other students of the team. We compare our results in the muon channel with the outcomes of the other two single lepton channels in the last chapter of this document. The final results are interpreted under the R-parity conserving MSSM, focusing on compressed slepton

democratic scenario (where all three sleptons are light and with similar mass). We consider different mass gap cases ($\Delta m = m_{\tilde{\chi}_1^\pm} - m_{\tilde{\chi}_1^0}$) ranging between 1 GeV and 50 GeV.

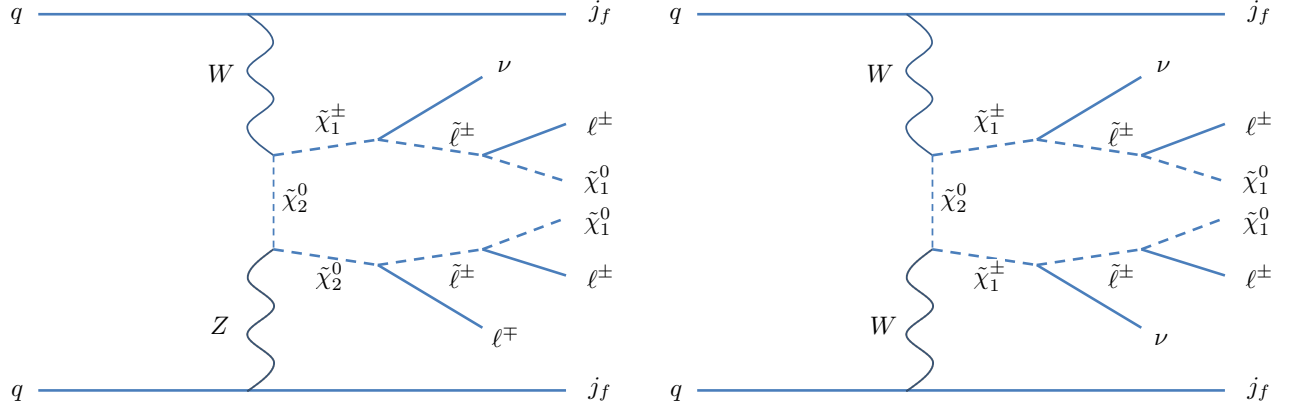


Figure 6.1: Diagrams of chargino-neutralino (left) and chargino-chargino (right) pair production through vector-boson fusion followed by their decays to leptons and $\tilde{\chi}_1^0$.

6.1 Strategy

The searches of SUSY using VBF tagging are similar to the searches that use cascade decays, such as colored SUSY searches, in the sense that high p_T jets ($p_T > 60$ GeV) are used to reduce the backgrounds. However, unlike the color searches, the production of squark and gluinos with subsequent decay chains is not required. Another important difference is the VBF topology of the jets: large pseudo-rapidity gap between them, produced in opposite hemispheres in the forward direction, and large dijet invariant mass. The production cross sections for electroweakinos is small (approximately 1–100 fb) and depends on the different mass values for each chargino and neutralino. Since the latter are mixed states of gauginos and higgsinos, it means that the mass values depend on the mixing matrix for each case. The identification of the VBF jets allows a robust discrimination between signal and SM backgrounds and acts as a counter-effect to the small electroweak production cross sections.

The whole set of selections is divided into two parts.

1. The set of central selections: composed of a MET requirement and the identification and selection of leptons.
2. The VBF selection: requiring at least one pair of high p_T jets with large pseudorapidity separation ($\Delta\eta$), in opposite hemispheres ($\eta_1 \times \eta_2 < 0$), and with large invariant mass (m_{jj}).

One distinctive characteristic of the signal is that the produced lepton in the final state (in our case, a muon) has a lower p_T spectrum than the backgrounds as it is depicted in Figure 6.2. This is because for our SUSY scenario, the LSP mass is slightly lower than the masses of other charginos and neutralinos (compressed spectra), which means that the decay products of the SUSY particles

produced would have low p_T . Based on the final states, the main SM backgrounds are $t\bar{t}$ and W+jets (which account for about 90% of the total background), with small contributions from other sources such as single top, dibosons, DY+jets and QCD. The VBF requirements reduce the background rates by a factor of $\sim 10^2$ - 10^4 , depending on the background origin. In addition to the VBF conditions, a relatively high E_T^{miss} is imposed due to the presence of the LSP, which leaves the detector undetected. The determination and optimization of the selections is discussed in Section 6.4. We estimate the remaining background yields after all requirements by the use of Monte-Carlo (MC) simulations. We use control regions (CR), with orthogonal requirements to the signal selection, to get a sample of real events from proton-proton collisions, with one background type dominating (with an expected purity greater than 60%). These CR's help us to verify how well the MC simulation describes the real collision data.

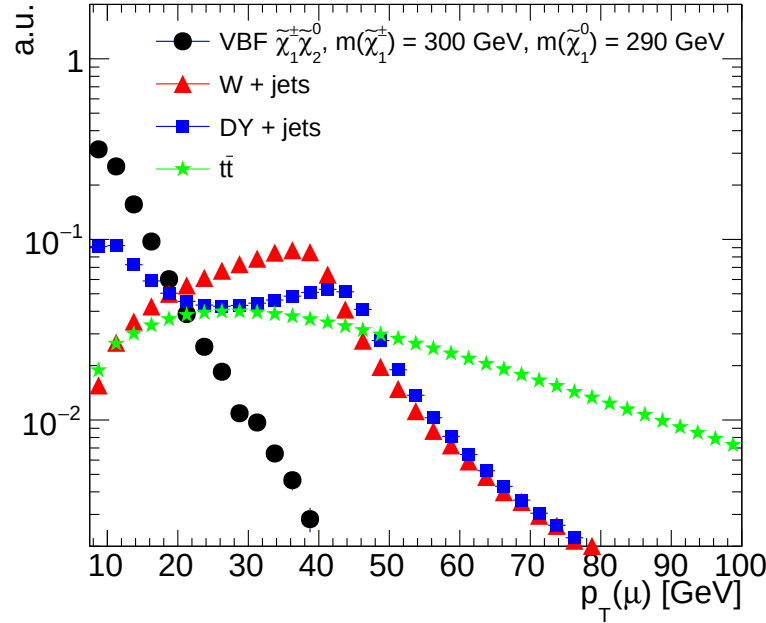


Figure 6.2: $p_T(\mu)$ distributions normalized to unity for three backgrounds: W+jets, DY+jets and $t\bar{t}$ and one VBF signal sample $\tilde{\chi}_1^\pm \tilde{\chi}_2^0$ with $m(\tilde{\chi}_1^\pm) = 300$ GeV and $m(\tilde{\chi}_1^0) = 290$ GeV. The distributions do not have any kinematic cut applied.

6.1.1 Correcting MC mismodeling

Our Signal Region is defined (SR) as the result of applying all the kinematic selections (explained in detail in the following sections) to the MC and data samples. The remaining MC yields in the SR, after all signal requirements, are compared to real proton-proton data with same requirements to look for an excess of events over background that would be the indication of new physics beyond SM. However, we have to verify that the MC simulates correctly the main two conditions we require: Central selections and VBF requirements. We perform this verification by defining two control regions (CR1 and CR2), one to check MC central selection mismodeling and the second one for

VBF requirements mismodeling. We first compare the distributions of different kinematic variables between data and MC in the Control regions. If the shapes are the same it is easier to correct the MC with just one normalization (scale factor) instead of doing a bin-by-bin correction. To avoid any correlation between the two types of requirements, the two control regions are necessarily chosen to have orthogonal conditions. These control regions are always orthogonal to the SR with the goal of obtaining two Scale Factors (SF1 and SF2) that are then applied to the MC yields to correct for any mismodeling. Note that no mismodeling would mean that the scale factors we obtain for each control region would be equal to unity. The idea is that these control regions should have a high purity of the background in consideration, in other words, each control region is an enriched sample of the background under study, but also, it is essential to have a negligible contribution from signal (to avoid any signal contamination). This is the reason why the control regions are defined as orthogonal to the SR. Figure 6.3 describes how we define the control regions. CR1 is used

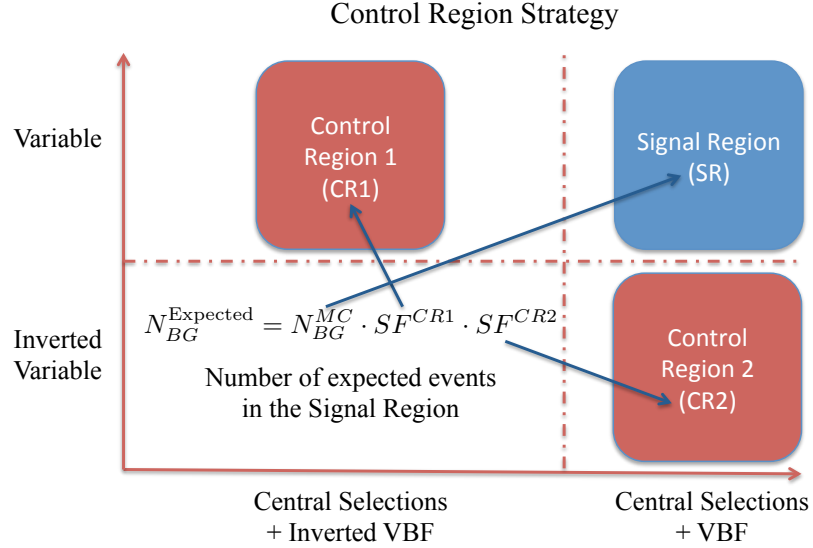


Figure 6.3: Diagram of the Control Regions (CR1 and CR2) used to obtain scale factors to correct MC mismodeling. The idea is to choose always orthogonal regions to the signal region.

to estimate a scale factor (SF1) for the central selections, that means, all the kinematic selections without the VBF requirement. Therefore it is defined by using inverted VBF selections. In principle, this inverted selection is fulfilled when two high- p_T jets are not available, or when the jets do not meet the VBF conditions as shown in Figure 6.4. CR2 is used to estimate a scale factor (SF2) for the VBF selections, that means, for CR2 it is mandatory to include the VBF selections, but at the same time it has to be orthogonal to CR1 and to the SR. To achieve this, we choose a “variable” that enriches the background sample, but at the same time, the inverted requirement on that variable allows to arrive at the signal region. To provide an example, for the $t\bar{t}$ background, we pick the number of b-jets as the variable to enrich the $t\bar{t}$ sample. In our signal region we require to have 0 b-jets, but for the control region, we demand to have exactly 1 b-jet. These conditions guarantee that we are orthogonal to the signal region, but, at the same time, we have an enriched sample of $t\bar{t}$ events. Section 6.5 provides more details.

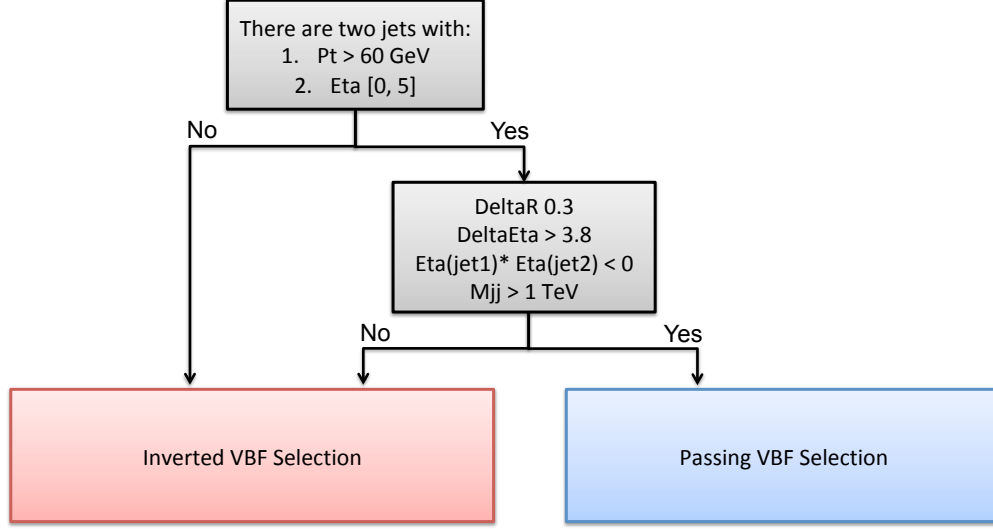


Figure 6.4: Diagram with the two conditions of the inverted VBF selection.

We use the following equation to estimate the surviving background yields in the SR:

$$N_{\text{BG}}^{\text{Data}} = N_{\text{BG}}^{\text{MC}}(\text{SR cuts}) \cdot SF_{\text{No VBF cuts}}^{\text{CR1}} \cdot SF_{\text{with VBF cuts}}^{\text{CR2}} \quad (6.1)$$

where $N_{\text{BG}}^{\text{Data}}$ is the predicted background yield in the SR, $N_{\text{BG}}^{\text{MC}}(\text{SR cuts})$ the predicted yield by MC simulation with the entire signal selections. $SF_{\text{No VBF cuts}}^{\text{CR1}}$ is the data-to-simulation scale factor for the central selections as determined from a background enriched control sample CR1 and given by the following equation:

$$SF_{\text{No VBF cuts}}^{\text{CR1}} = \frac{N^{\text{CR1}}(\text{Data}) - N^{\text{CR1}}(\text{other BGs})}{N^{\text{CR1}}(\text{BG studied})} \quad (6.2)$$

where $N^{\text{CR1}}(\text{Data})$ is the number of data events passing the CR1 selection, $N^{\text{CR1}}(\text{BG studied})$ is the number of MC events passing the CR1 selection for the background sample under study (for example $t\bar{t}$ events), and $N^{\text{CR1}}(\text{other BGs})$ is the sum of all the backgrounds different to the background under study (in our example, different to $t\bar{t}$ events). $SF_{\text{with VBF cuts}}^{\text{CR2}}$ is the data-to-simulation scale factor for the VBF selections as determined in a different background enriched control sample CR2 and given by the following equation:

$$SF_{\text{with VBF cuts}}^{\text{CR2}} = \frac{N^{\text{CR2}}(\text{Data}) - N^{\text{CR2}}(\text{other BGs})}{N^{\text{CR2}}(\text{BG studied}) \cdot SF_{\text{No VBF cuts}}^{\text{CR1}}} \quad (6.3)$$

where $N^{\text{CR2}}(\text{Data})$ is the number of data events passing the CR2 selection, $N^{\text{CR2}}(\text{BG studied})$ is the number of MC events passing the CR2 selection for the background sample under study, $N^{\text{CR2}}(\text{other BGs})$ is the sum of all the backgrounds different to the background under study and $SF_{\text{No VBF cuts}}^{\text{CR1}}$ is the SF calculated for CR1 defined in Equation 6.2. Figure 6.3 indicates the correlation of SF's and their corresponding CR's. It is important to note that the data-to-simulation scale factor for the central selections, $SF_{\text{No VBF cuts}}^{\text{CR1}}$, includes the corrections for everything related

to central selections: object identification and isolation, MET requirement, vetoes on other leptons, etc. Another important requirement for the control regions is to assure that $SF_{\text{No VBF cuts}}^{\text{CR1}}$ ($SF_{\text{with VBF cuts}}^{\text{CR2}}$) is not biased by the definition of CR1 (CR2). Additional regions are defined in order to perform closure tests on this method. More details are given in Sections 6.5.1 and 6.6.1.

6.2 Data and Monte Carlo Samples

All the data used in the present study comes from proton-proton collisions carried out at $\sqrt{s}$ of 13 TeV during 2016 recorded by the CMS experiment. A MET Primary Data Set is used for most of the analysis. A single muon data set was used for trigger studies and for a closure test of $t\bar{t}$ background. The MET data set is characterized by filtering events based on a MET requirement without demanding any other additional condition. This is important because we do not want to have any bias on the p_T of the lepton. Table 6.1 shows the complete set of data used. The collision data sample corresponds to an integrated luminosity of 35.86 fb^{-1} . Table 6.2 shows the whole set of MC samples used, together with their respective cross-sections. This data set is part of the official MC production carried out by the CMS experiment during 2016 (Summer miniAODv2 2016). Pythia8 and Madgraph generators, that simulate processes to leading order (LO), were used in the MC production. However, to obtain the background estimates, we use next-to-leading (NLO) and next-to-next-to-leading order (NNLO) cross sections. We normalize all the MC samples by using the cross section values in Table 6.2. Table 6.3 shows all the signal samples used. The central condition for generating the samples is the production of charginos and neutralinos through VBF processes, removing the production of QCD vertices. There is a correction performed to the MC samples since the pileup distribution of MC is different from the pileup distribution of data. The way to correct the MC is to assign a weight to each bin of the MC pileup distribution to match the data. We determine the probability of obtaining ‘n’ interactions in MC ($P_{MC}(n)$) and data ($P_{data}(n)$), to calculate the weight that has to be applied to every bin of the MC pileup distribution, which corresponds to the fraction:

$$w_{PU}(n) = \frac{P_{data}(n)}{P_{MC}(n)} \quad (6.4)$$

The previous result would serve to multiply the bin ‘n’ of the MC pileup distribution corresponding to an ‘n’ number of interactions.

6.3 Triggers

A specific VBF trigger was developed at CMS experiment for Run 2 of the LHC. This trigger includes selection of two VBF jets, E_T^{miss} and muons with a p_T above threshold (8 GeV). Since we are selecting soft leptons in our signal topology (a soft muon as shown in Figure 6.2), we pick a trigger that do not make any requirement on the p_T of the leptons, to avoid any bias coming from the trigger selection. In this way, collision data events that have passed through a “MET trigger” (HLT_PFMETNoMu120_PFMHTNoMu120_IDTight) are selected. The different sections on the trigger name have the following meaning:

- HLT: This filter is selected in the High Level Trigger. The trigger has the L1 seeds listed below. Please note that all of them are selecting E_T^{miss} greater than a certain threshold.

Table 6.1: Collision Data Samples

Physics Sample	Official CMS Datasets
Run 2016B SingleMu Run2016B-03Feb2017	<i>/SingleMuon/Run2016B-03Feb2017_ver2-v2/MINIAOD</i>
Run 2016C SingleMu Run2016C-03Feb2017	<i>/SingleMuon/Run2016C-03Feb2017-v1/MINIAOD</i>
Run 2016D SingleMu Run2016D-03Feb2017	<i>/SingleMuon/Run2016D-03Feb2017-v1/MINIAOD</i>
Run 2016E SingleMu Run2016E-03Feb2017	<i>/SingleMuon/Run2016E-03Feb2017-v1/MINIAOD</i>
Run 2016F SingleMu Run2016F-03Feb2017	<i>/SingleMuon/Run2016F-03Feb2017-v1/MINIAOD</i>
Run 2016G SingleMu Run2016G-03Feb2017	<i>/SingleMuon/Run2016G-03Feb2017-v1/MINIAOD</i>
Run 2016Hv2 SingleMu Run2016H-03Feb2017	<i>/SingleMuon/Run2016H-03Feb2017_ver2-v1/MINIAOD</i>
Run 2016Hv3 SingleMu Run2016H-03Feb2017	<i>/SingleMuon/Run2016H-03Feb2017_ver3-v1/MINIAOD</i>
Run 2016B Met Run2016B-03Feb2017	<i>/Met/Run2016B-03Feb2017-v3/MINIAOD</i>
Run 2016C Met Run2016C-03Feb2017	<i>/Met/Run2016C-03Feb2017-v1/MINIAOD</i>
Run 2016D Met Run2016D-03Feb2017	<i>/Met/Run2016D-03Feb2017-v1/MINIAOD</i>
Run 2016E Met Run2016E-03Feb2017	<i>/Met/Run2016E-03Feb2017-v1/MINIAOD</i>
Run 2016F Met Run2016F-03Feb2017	<i>/Met/Run2016F-03Feb2017-v1/MINIAOD</i>
Run 2016G Met Run2016G-03Feb2017	<i>/Met/Run2016G-03Feb2017-v1/MINIAOD</i>
Run 2016Hv2 Met Run2016H-03Feb2017	<i>/Met/Run2016H-03Feb2017_ver2-v1/MINIAOD</i>
Run 2016Hv3 Met Run2016H-03Feb2017	<i>/Met/Run2016H-03Feb2017_ver3-v1/MINIAOD</i>

- L1_ETM50
- L1_ETM60
- L1_ETM70
- L1_ETM75
- L1_ETM80
- L1_ETM85
- L1_ETM90
- L1_ETM95
- L1_ETM100
- L1_ETM120
- L1_ETM75_Jet60_dPhi_Min0p4
- L1_DoubleJetC60_ETM60
- PFMETNoMu120: This trigger section implies that events that satisfy the condition of E_T^{miss} greater than 120 GeV are selected. During the E_T^{miss} reconstruction all the events that have muons reconstructed with the PF algorithm are rejected due to the low deposition of energy left by muons in the detector.
- PFMHTNoMu120: This trigger section implies that events with reconstructed jets with a requirement in the MHT variable greater than 120 GeV are selected. During the reconstruction of these jets, all the events that have muons reconstructed with the PF algorithm as part of the jets, are rejected.

Table 6.2: MC Samples and their respective cross sections.

Process	cross-section [pb]	Official CMS Datasets (MINIAODSIM)
$Z \rightarrow \mu\mu$ HT binned LO	175.3×1.17781 (NNLO)	/DYJetsToLL_M-50_HT-70to100_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	147.4×1.17781 (NNLO)	/DYJetsToLL_M-50_HT-100to200_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	147.4×1.17781 (NNLO)	/DYJetsToLL_M-50_HT-100to200_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext1-v1/MINIAODSIM
	40.99×1.17781 (NNLO)	/DYJetsToLL_M-50_HT-200to400_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	40.99×1.17781 (NNLO)	/DYJetsToLL_M-50_HT-200to400_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext1-v1/MINIAODSIM
	5.678×1.17781 (NNLO)	/DYJetsToLL_M-50_HT-400to600_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	5.678×1.17781 (NNLO)	/DYJetsToLL_M-50_HT-400to600_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext1-v1/MINIAODSIM
	1.363×1.17781 (NNLO)	/DYJetsToLL_M-50_HT-600to800_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v2/MINIAODSIM
	0.6759×1.17781 (NNLO)	/DYJetsToLL_M-50_HT-800to1200_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	0.116×1.17781 (NNLO)	/DYJetsToLL_M-50_HT-1200to2500_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	0.002592×1.17781 (NNLO)	/DYJetsToLL_M-50_HT-2500toInf_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
EWK $Z \rightarrow \mu\mu$	3.998 (LO)	/EWKZ2Jets_ZToLL_M-50_13TeV-madgraph-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext2-v1/MINIAODSIM
$W + \text{jets}$ HT binned LO	1319×1.21378 (NNLO)	/WJetsToLNu_HT-70To100_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	1345×1.21378 (NNLO)	/WJetsToLNu_HT-100To200_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	1345×1.21378 (NNLO)	/WJetsToLNu_HT-100To200_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext1-v1/MINIAODSIM
	359.7×1.21378 (NNLO)	/WJetsToLNu_HT-200To400_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	359.7×1.21378 (NNLO)	/WJetsToLNu_HT-200To400_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext1-v1/MINIAODSIM
	48.91×1.21378 (NNLO)	/WJetsToLNu_HT-400To600_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	48.91×1.21378 (NNLO)	/WJetsToLNu_HT-400To600_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext1-v1/MINIAODSIM
	12.05×1.21378 (NNLO)	/WJetsToLNu_HT-600To800_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	12.05×1.21378 (NNLO)	/WJetsToLNu_HT-600To800_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext1-v1/MINIAODSIM
	5.501×1.21378 (NNLO)	/WJetsToLNu_HT-800To1200_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	5.501×1.21378 (NNLO)	/WJetsToLNu_HT-800To1200_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext1-v1/MINIAODSIM
	1.329×1.21378 (NNLO)	/WJetsToLNu_HT-1200To2500_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	1.329×1.21378 (NNLO)	/WJetsToLNu_HT-1200To2500_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext1-v1/MINIAODSIM
	0.03216×1.21378 (NNLO)	/WJetsToLNu_HT-2500ToInf_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	0.03216×1.21378 (NNLO)	/WJetsToLNu_HT-2500ToInf_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext1-v1/MINIAODSIM
EWK $W + \text{jets}$	25.81 (LO) 20.35 (LO)	/EWKWPplus2Jets_WToLNu_M-50_13TeV-madgraph-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext2-v1/MINIAODSIM /EWKWPminus2Jets_WToLNu_M-50_13TeV-madgraph-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext2-v1/MINIAODSIM
QCD	24670000 (LO)	/QCD_HT50To100_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	28080000 (LO)	/QCD_HT100To200_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	1712000 (LO)	/QCD_HT200To300_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	347500 (LO)	/QCD_HT300To500_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	32100 (LO)	/QCD_HT500To700_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	6823 (LO)	/QCD_HT700To1000_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	1208 (LO)	/QCD_HT1000To1500_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	120 (LO)	/QCD_HT1500To2000_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	25.26 (LO)	/QCD_HT2000ToInf_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
$Z \rightarrow \nu\nu + \text{jets}$ HT binned LO	280.35 (LO)	/ZJetsToNuNu_HT-100To200_13TeV-madgraph/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	77.67 (LO)	/ZJetsToNuNu_HT-200To400_13TeV-madgraph/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	10.73 (LO)	/ZJetsToNuNu_HT-400To600_13TeV-madgraph/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	0.853 (LO)	/ZJetsToNuNu_HT-600To800_13TeV-madgraph/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	0.3942 (LO)	/ZJetsToNuNu_HT-800To1200_13TeV-madgraph/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	0.0974 (LO)	/ZJetsToNuNu_HT-1200To2500_13TeV-madgraph/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	0.002308 (LO)	/ZJetsToNuNu_HT-2500ToInf_13TeV-madgraph/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
EWK $Z \rightarrow \nu\nu$	3.998 (LO)	/EWKZ2Jets_ZToNuNu_13TeV-madgraph-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
Top quark	831.76 (NNLO)	/TT_TuneCUETP8M2T4_13TeV-powheg-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
	35.6 (LO)	/ST_tW_antitop_5f_inclusiveDecays_13TeV-powheg-pythia8_TuneCUETP8M1/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext1-v1
	35.6 (LO)	/ST_tW_top_5f_inclusiveDecays_13TeV-powheg-pythia8_TuneCUETP8M1/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6_ext1-v1
	136.02 (LO)	/ST_t-channel_top_4f_inclusiveDecays_13TeV-powhegV2-madspin-pythia8_TuneCUETP8M1/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1
	80.95 (LO)	/ST_t-channel_antitop_4f_inclusiveDecays_13TeV-powhegV2-madspin-pythia8_TuneCUETP8M1/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1
WW	63.21 (LO)	/WW_TuneCUETP8M1_13TeV-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
WZ	22.82 (LO)	/WZ_TuneCUETP8M1_13TeV-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
ZZ	10.32 (LO)	/ZZ_TuneCUETP8M1_13TeV-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
VBS		/WWJJToLNuLNu_EWK_13TeV-madgraph-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM /WZJJ_EWK_13TeV-madgraph-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM /ZZJJToLL_EWK_13TeV-madgraph-pythia8/RunIISummer16MiniAODv2-PUMoriond17_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM
VBF Signal		/VBF-C1n2_leptonicDecays_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUSummer16Fast_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM /VBF-C1n2_WZ_TuneCUETP8M1_13TeV-madgraphMLM-pythia8/RunIISummer16MiniAODv2-PUSummer16Fast_80X_mcRun2_asymptotic_2016_TracheIV_v6-v1/MINIAODSIM

Table 6.3: Signal mass points and their respective cross sections.

Process	xs [fb]	Process	xs [fb]	Process	xs [fb]	Process	xs [fb]	Process	xs [fb]	Process	xs [fb]
VBF-C1N2 leptonic 100 50	286.3	VBF-C1N2 WZ 100 50	287.8	VBF-C1N2 leptonic 250 200	16.0	VBF-C1N2 WZ 250 200	15.9	VBF-C1N2 leptonic 400 350	3.1	VBF-C1N2 WZ 400 350	3.1
VBF-C1N2 leptonic 100 55		VBF-C1N2 WZ 100 55		VBF-C1N2 leptonic 250 205		VBF-C1N2 WZ 250 205		VBF-C1N2 leptonic 400 355		VBF-C1N2 WZ 400 355	
VBF-C1N2 leptonic 100 60		VBF-C1N2 WZ 100 60		VBF-C1N2 leptonic 250 210		VBF-C1N2 WZ 250 210		VBF-C1N2 leptonic 400 360		VBF-C1N2 WZ 400 360	
VBF-C1N2 leptonic 100 65		VBF-C1N2 WZ 100 65		VBF-C1N2 leptonic 250 215		VBF-C1N2 WZ 250 215		VBF-C1N2 leptonic 400 365		VBF-C1N2 WZ 400 365	
VBF-C1N2 leptonic 100 70		VBF-C1N2 WZ 100 70		VBF-C1N2 leptonic 250 220		VBF-C1N2 WZ 250 220		VBF-C1N2 leptonic 400 370		VBF-C1N2 WZ 400 370	
VBF-C1N2 leptonic 100 75		VBF-C1N2 WZ 100 75		VBF-C1N2 leptonic 250 225		VBF-C1N2 WZ 250 225		VBF-C1N2 leptonic 400 375		VBF-C1N2 WZ 400 375	
VBF-C1N2 leptonic 100 80		VBF-C1N2 WZ 100 80		VBF-C1N2 leptonic 250 230		VBF-C1N2 WZ 250 230		VBF-C1N2 leptonic 400 380		VBF-C1N2 WZ 400 380	
VBF-C1N2 leptonic 100 85		VBF-C1N2 WZ 100 85		VBF-C1N2 leptonic 250 235		VBF-C1N2 WZ 250 235		VBF-C1N2 leptonic 400 385		VBF-C1N2 WZ 400 385	
VBF-C1N2 leptonic 100 90		VBF-C1N2 WZ 100 90		VBF-C1N2 leptonic 250 240		VBF-C1N2 WZ 250 240		VBF-C1N2 leptonic 400 390		VBF-C1N2 WZ 400 390	
VBF-C1N2 leptonic 100 95		VBF-C1N2 WZ 100 95		VBF-C1N2 leptonic 250 245		VBF-C1N2 WZ 250 245		VBF-C1N2 leptonic 400 395		VBF-C1N2 WZ 400 395	
VBF-C1N2 leptonic 100 98		VBF-C1N2 WZ 100 98		VBF-C1N2 leptonic 250 248		VBF-C1N2 WZ 250 248		VBF-C1N2 leptonic 400 398		VBF-C1N2 WZ 400 398	
VBF-C1N2 leptonic 100 99		VBF-C1N2 WZ 100 99		VBF-C1N2 leptonic 250 249		VBF-C1N2 WZ 250 249		VBF-C1N2 leptonic 400 399		VBF-C1N2 WZ 400 399	
VBF-C1N2 leptonic 150 100	80.5	VBF-C1N2 WZ 150 100	80.5	VBF-C1N2 leptonic 300 250	8.7	VBF-C1N2 WZ 300 250	8.6	VBF-C1N2 leptonic 450 400	2.0	VBF-C1N2 WZ 450 400	2.0
VBF-C1N2 leptonic 150 105		VBF-C1N2 WZ 150 105		VBF-C1N2 leptonic 300 255		VBF-C1N2 WZ 300 255		VBF-C1N2 leptonic 450 405		VBF-C1N2 WZ 450 405	
VBF-C1N2 leptonic 150 110		VBF-C1N2 WZ 150 110		VBF-C1N2 leptonic 300 260		VBF-C1N2 WZ 300 260		VBF-C1N2 leptonic 450 410		VBF-C1N2 WZ 450 410	
VBF-C1N2 leptonic 150 115		VBF-C1N2 WZ 150 115		VBF-C1N2 leptonic 300 265		VBF-C1N2 WZ 300 265		VBF-C1N2 leptonic 450 415		VBF-C1N2 WZ 450 415	
VBF-C1N2 leptonic 150 120		VBF-C1N2 WZ 150 120		VBF-C1N2 leptonic 300 270		VBF-C1N2 WZ 300 270		VBF-C1N2 leptonic 450 420		VBF-C1N2 WZ 450 420	
VBF-C1N2 leptonic 150 125		VBF-C1N2 WZ 150 125		VBF-C1N2 leptonic 300 275		VBF-C1N2 WZ 300 275		VBF-C1N2 leptonic 450 425		VBF-C1N2 WZ 450 425	
VBF-C1N2 leptonic 150 130		VBF-C1N2 WZ 150 130		VBF-C1N2 leptonic 300 280		VBF-C1N2 WZ 300 280		VBF-C1N2 leptonic 450 430		VBF-C1N2 WZ 450 430	
VBF-C1N2 leptonic 150 135		VBF-C1N2 WZ 150 135		VBF-C1N2 leptonic 300 285		VBF-C1N2 WZ 300 285		VBF-C1N2 leptonic 450 435		VBF-C1N2 WZ 450 435	
VBF-C1N2 leptonic 150 140		VBF-C1N2 WZ 150 140		VBF-C1N2 leptonic 300 290		VBF-C1N2 WZ 300 290		VBF-C1N2 leptonic 450 440		VBF-C1N2 WZ 450 440	
VBF-C1N2 leptonic 150 145		VBF-C1N2 WZ 150 145		VBF-C1N2 leptonic 300 295		VBF-C1N2 WZ 300 295		VBF-C1N2 leptonic 450 445		VBF-C1N2 WZ 450 445	
VBF-C1N2 leptonic 150 148		VBF-C1N2 WZ 150 148		VBF-C1N2 leptonic 300 298		VBF-C1N2 WZ 300 298		VBF-C1N2 leptonic 450 448		VBF-C1N2 WZ 450 448	
VBF-C1N2 leptonic 150 149		VBF-C1N2 WZ 150 149		VBF-C1N2 leptonic 300 299		VBF-C1N2 WZ 300 299		VBF-C1N2 leptonic 450 449		VBF-C1N2 WZ 450 449	
VBF-C1N2 leptonic 200 150	32.6	VBF-C1N2 WZ 200 150	32.6	VBF-C1N2 leptonic 350 300	5.1	VBF-C1N2 WZ 350 300	5.0	VBF-C1N2 leptonic 500 450	1.4	VBF-C1N2 WZ 500 450	1.4
VBF-C1N2 leptonic 200 155		VBF-C1N2 WZ 200 155		VBF-C1N2 leptonic 350 305		VBF-C1N2 WZ 350 305		VBF-C1N2 leptonic 500 455		VBF-C1N2 WZ 500 455	
VBF-C1N2 leptonic 200 160		VBF-C1N2 WZ 200 160		VBF-C1N2 leptonic 350 310		VBF-C1N2 WZ 350 310		VBF-C1N2 leptonic 500 460		VBF-C1N2 WZ 500 460	
VBF-C1N2 leptonic 200 165		VBF-C1N2 WZ 200 165		VBF-C1N2 leptonic 350 315		VBF-C1N2 WZ 350 315		VBF-C1N2 leptonic 500 465		VBF-C1N2 WZ 500 465	
VBF-C1N2 leptonic 200 170		VBF-C1N2 WZ 200 170		VBF-C1N2 leptonic 350 320		VBF-C1N2 WZ 350 320		VBF-C1N2 leptonic 500 470		VBF-C1N2 WZ 500 470	
VBF-C1N2 leptonic 200 175		VBF-C1N2 WZ 200 175		VBF-C1N2 leptonic 350 325		VBF-C1N2 WZ 350 325		VBF-C1N2 leptonic 500 475		VBF-C1N2 WZ 500 475	
VBF-C1N2 leptonic 200 180		VBF-C1N2 WZ 200 180		VBF-C1N2 leptonic 350 330		VBF-C1N2 WZ 350 330		VBF-C1N2 leptonic 500 480		VBF-C1N2 WZ 500 480	
VBF-C1N2 leptonic 200 185		VBF-C1N2 WZ 200 185		VBF-C1N2 leptonic 350 335		VBF-C1N2 WZ 350 335		VBF-C1N2 leptonic 500 485		VBF-C1N2 WZ 500 485	
VBF-C1N2 leptonic 200 190		VBF-C1N2 WZ 200 190		VBF-C1N2 leptonic 350 340		VBF-C1N2 WZ 350 340		VBF-C1N2 leptonic 500 490		VBF-C1N2 WZ 500 490	
VBF-C1N2 leptonic 200 195		VBF-C1N2 WZ 200 195		VBF-C1N2 leptonic 350 345		VBF-C1N2 WZ 350 345		VBF-C1N2 leptonic 500 495		VBF-C1N2 WZ 500 495	
VBF-C1N2 leptonic 200 198		VBF-C1N2 WZ 200 198		VBF-C1N2 leptonic 350 348		VBF-C1N2 WZ 350 348		VBF-C1N2 leptonic 500 498		VBF-C1N2 WZ 500 498	
VBF-C1N2 leptonic 200 199		VBF-C1N2 WZ 200 199		VBF-C1N2 leptonic 350 349		VBF-C1N2 WZ 350 349		VBF-C1N2 leptonic 500 499		VBF-C1N2 WZ 500 499	

To study the MET trigger selected, we use a sample that contains W+jets events. To ensure that we have an enriched sample of W+jets events we apply the following requirements:

- We use a muon dataset to which a single-muon trigger (HLT_IsoMu24_Eta2p1_v) is applied. Our goal with this condition is to obtain events with a muon produced by a W decay ($W \rightarrow \mu\nu$).
- MuonTight ID, $p_T > 30$ GeV, $|\eta| < 2.1$. To get well-reconstructed muons.
- The VBF selections for the jets, to have events in the final state with the same topology of our signal.
- A low E_T^{miss} condition (> 50 GeV) to study the efficiency of the trigger as a function of E_T^{miss} .

Once the above requirements are applied, we define our numerator including those events that pass through the MET trigger. We define the efficiency as follows:

$$\epsilon = \frac{\text{The denominator \& events firing HLT_PFMETNoMu120_PFMHTNoMu120_IDTight}}{\text{Events firing HLT_IsoMu24_Eta2p1_v \& Tight ID \& } p_T > 30\text{GeV \& } |\eta| < 2.1 \text{ \& VBF selections}} \quad (6.5)$$

The ‘&’ symbol in Equation 6.5 is acting as an AND operator. Figure 6.5 shows the trigger efficiency for the HLT_PFMETNoMu120_PFMHTNoMu120_IDTight trigger as a function of E_T^{miss} for data and MC. The left figure corresponds to the case where both jets are central ($|\eta| < 3$), while the right figure corresponds to the condition with one jet central and the other jet forward ($3 < |\eta| < 5$). The figure includes a sample containing VBF jets (VBF W process) to study the performance of the trigger with an enriched VBF sample. Figure 6.5 shows that all the efficiency curves reach $\sim 100\%$ plateaus for values close to E_T^{miss} of 250 GeV. For lower E_T^{miss} values, the efficiency curves have different behaviour. In particular, there is a difference between the data sample (single μ) and the other two MC samples. The choice of $E_T^{\text{miss}} > 250$ GeV avoids this difference. However, it should be noted that in all the results of this analysis, the trigger is applied only to the data and not to the MC. This is because the trigger is not well modeled in the MC samples. Instead, the same kinematic selections for data and MC are applied.

6.4 Event Selection

6.4.1 Lepton selection and b-jet requirement

The selection of low p_T muons (between 8 and 40 GeV) is motivated by the type of SUSY signals we are searching for. Figure 6.2 shows that for a signal benchmark point with a mass gap $\Delta m = m(\tilde{\chi}_1^\pm) - m(\tilde{\chi}_1^0) = 10$ GeV the signal is more pronounced than the backgrounds in the region 8 GeV to 20 GeV, therefore, this region is where we would expect the highest significance. The p_T window selected allows to include other SUSY scenarios with higher mass gaps. For the case of our interest the LSP mass is slightly lower than the masses of other charginos and neutralinos, which means that the decay products of the SUSY particles produced would have low p_T , which makes it difficult to reconstruct and identify multiple leptons, consequently, we focus on events that have exactly one well-identified soft muon. To ensure that we have exactly one muon in the final state, a veto is applied on well-reconstructed taus and electrons. In addition, all leptons are required to have pseudorapidity $|\eta| < 2.1$ in order to ensure that they are reconstructed fully within the tracking system. Muon candidates are also required to pass the reconstruction and identification criteria

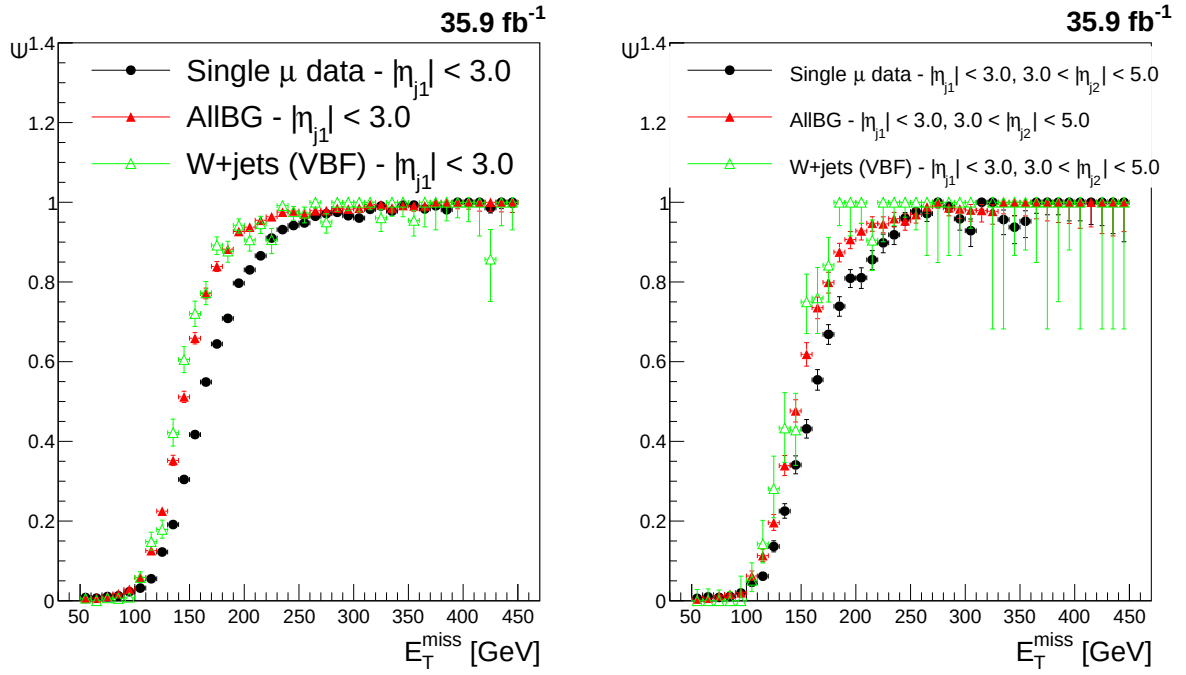


Figure 6.5: Trigger efficiency for the HLT_PFMETNoMu120_PFMHTNoMu120_IDTight trigger. Left: trigger efficiency as a function of offline E_T^{miss} , for events where both jets have $|\eta| < 3$. Right: trigger efficiency as a function of offline E_T^{miss} , where one of the jets has $|\eta| < 3$ and the other $3 < |\eta| < 5$.

described in Section 5.5.5. In addition, muons have to be isolated to guarantee they do not come from hadronic decays. We require that the events do not have any jet identified as a b-quark jet (b-jet) to reduce the background from $t\bar{t}$ production, that indeed is the source with the highest background contribution to our signal region (more details are given in Section 6.5).

6.4.2 E_T^{miss} requirement

The high E_T^{miss} is required mainly motivated by the presence of LSP's in the final state of the decay chain of SUSY particles. Since we are assuming R-parity conserving models, all the SUSY particles have to be produced in pairs, each having a cascade decay ending with the production of a LSP, therefore we expect to have two LSP's in the final state, which do not interact with the detector and its presence is only indirectly inferred through E_T^{miss} . Since the LSP's are expected to be heavier than neutrinos, we would expect a higher E_T^{miss} distribution for a SUSY signal than for neutrinos produced by SM backgrounds. For events where the neutrinos have high kinetic energy while the LSP's are produced near at rest, the E_T^{miss} distributions would overlap (see Figure 6.6), and in this case we end up with a source of a non-reducible background. From Figure 6.6 one can also infer that a high E_T^{miss} requirement ($E_T^{\text{miss}} > 250$ GeV) suppresses drastically the backgrounds improving the signal significance, which is defined as:

$$S = \frac{s}{\sqrt{\sigma_s^2 + \sigma_B^2}} = \frac{s}{\sqrt{s + B}} \quad (6.6)$$

where s is the event yield from the signal and B is the sum of all background event yields. σ_s and σ_B are the uncertainties in signal and background events, respectively. Note that the signal significance, defined in this way can be understood as a measure of number of standard deviations that the signal stands over the background. Figure 6.7 shows that the best signal significance is achieved after 250 GeV, for the two signal samples:

- $m(\tilde{\chi}_1^\pm) = 200$ GeV, $m(\tilde{\tau}) = 160$ GeV, $m(\tilde{\chi}_1^0) = 150$ GeV
- $m(\tilde{\chi}_1^\pm) = 200$ GeV, $m(\tilde{\tau}) = 195$ GeV, $m(\tilde{\chi}_1^0) = 150$ GeV

In addition, Figure 6.8 shows the efficiency as function of E_T^{miss} cut for three backgrounds: W+jets, DY+jets and $t\bar{t}$ and one SUSY VBF signal sample with $m(\tilde{\chi}_1^\pm) = 300$ GeV and $m(\tilde{\chi}_1^0) = 299$ GeV. From this later plot one can conclude that the E_T^{miss} requirement of 250 GeV reduces in several orders of magnitude the background contributions (in particular, DY+jets), while the signal remains within acceptable statistics. Also this is an argument for not increasing the E_T^{miss} cut to higher levels, to avoid a decrease in statistics without any gain in significance.

6.4.3 $m_T(\mu, E_T^{\text{miss}})$ requirement

Our main motivation for applying an $m_T(\mu, E_T^{\text{miss}})$ condition is the suppression of the W+jets background. As shown in Figure 6.9, a requirement of $m_T > 110$ GeV guarantees that we are in the region beyond the jacobian peak of the W mass. Indeed, for values of m_T higher than 110 GeV, the shape of the signal begins to stand out over the normalized background shapes. In addition, the m_T requirement is also motivated by the signal significance, as shown in Figure 6.10, $m_T > 110$ GeV provides the best signal significance.

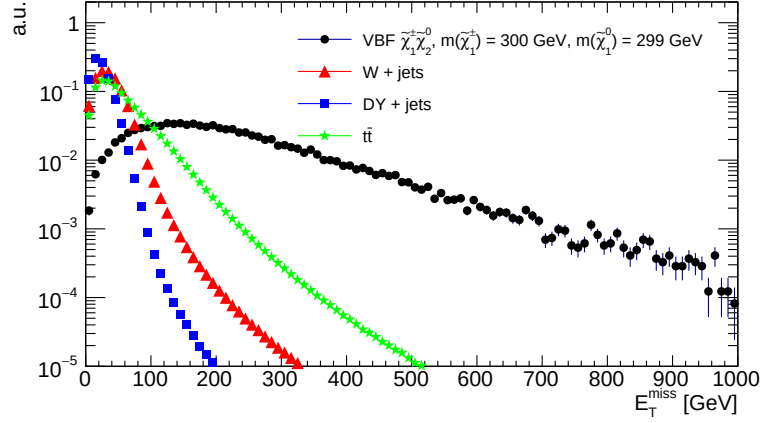


Figure 6.6: E_T^{miss} distributions normalized to unity for three backgrounds: W+jets, DY+jets and $t\bar{t}$ and one VBF signal sample $\tilde{\chi}_1^\pm \tilde{\chi}_2^0$ with $m(\tilde{\chi}_1^\pm) = 300$ GeV and $m(\tilde{\chi}_1^0) = 299$ GeV. The distributions do not have any kinematic requirement.

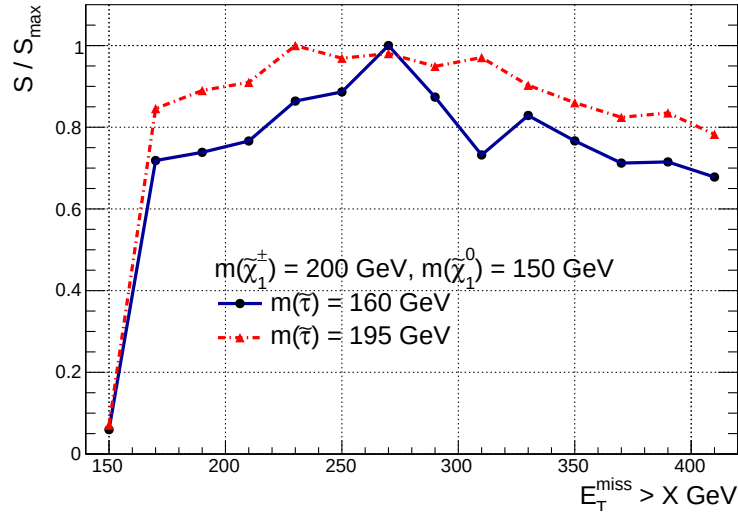


Figure 6.7: Signal significance $S = \frac{s}{\sqrt{s+B}}$ divided by the maximum significance S_{max} , as a function of E_T^{miss} cut values for two signal samples 1) $m(\tilde{\chi}_1^\pm) = 200$ GeV, $m(\tilde{\tau}) = 160$ GeV, $m(\tilde{\chi}_1^0) = 150$ GeV and 2) $m(\tilde{\chi}_1^\pm) = 200$ GeV, $m(\tilde{\tau}) = 195$ GeV, $m(\tilde{\chi}_1^0) = 150$ GeV. The best signal significance is achieved for $E_T^{\text{miss}} > 250$ GeV, it begins to decrease for higher E_T^{miss} requirements.

6.4.4 VBF selections

The VBF topology is characterized by the presence of two energetic jets in the forward direction, in opposite hemispheres, and with large dijet invariant mass. Figure 6.11 shows the distributions of the pseudorapidity variable η for all jets and the largest dijet invariant mass m_{jj} for the three main backgrounds added together ($t\bar{t}$, W+jets and DY+jets) compared to one SUSY VBF specific signal

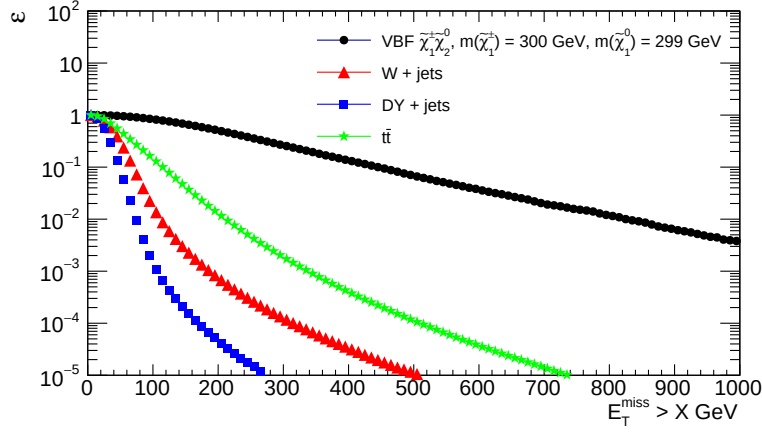


Figure 6.8: Signal and background efficiency as a function of E_T^{miss} cut.

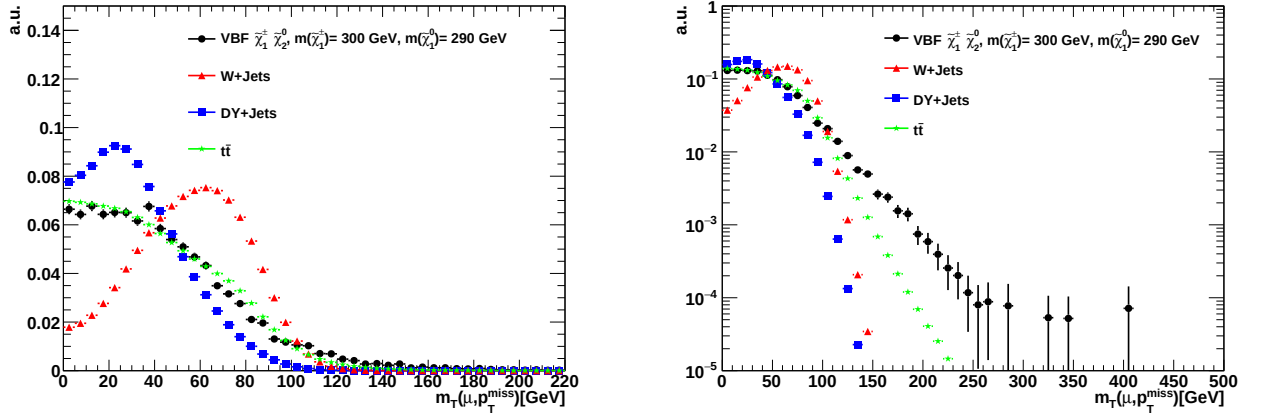


Figure 6.9: (Left) $m_T(\mu, p_T^{\text{miss}})$ distributions normalized to unity (without any kinematic cut applied) for three backgrounds: W+jets, DY+jets and $t\bar{t}$ and one SUSY VBF signal sample $\tilde{\chi}_1^\pm \tilde{\chi}_2^0$ with $m(\tilde{\chi}_1^\pm) = 300$ GeV and $m(\tilde{\chi}_2^0) = 290$ GeV. (Right) The previous distributions with logarithmic scale in the Y axis.

sample. According to these plots, the jets produced by the background events are mostly in the central region of the detector ($|\eta| \leq 1.8$) and have small dijet invariant masses, while the signal events are characterized by non-central jets (the two peaks in the upper plot of Figure 6.11 at $\eta \approx \pm 1.6$) with large dijet invariant masses (lower plot of Figure 6.11). The large m_{jj} is a consequence of the kinematics of the selected pair of jets in the event with a high pseudorapidity gap ($\Delta\eta_{jj}$) between them. The outgoing partons in VBF processes carry relatively large p_T since they must have enough momentum to radiate the SM vector bosons and keep moving in the forward direction. Therefore, the VBF selections are defined by requiring at least two jets, each with a p_T greater than 60 GeV and $|\eta| < 5.0$. Only jets separated from the leptons by $\Delta R = \sqrt{\Delta\phi^2 + \Delta\eta^2} > 0.3$ are considered. All jet candidates passing the above requirements and having $|\Delta\eta_{jj}| > 3.8$, $\eta_1 \cdot \eta_2 < 0$, and $m_{jj} > 1$

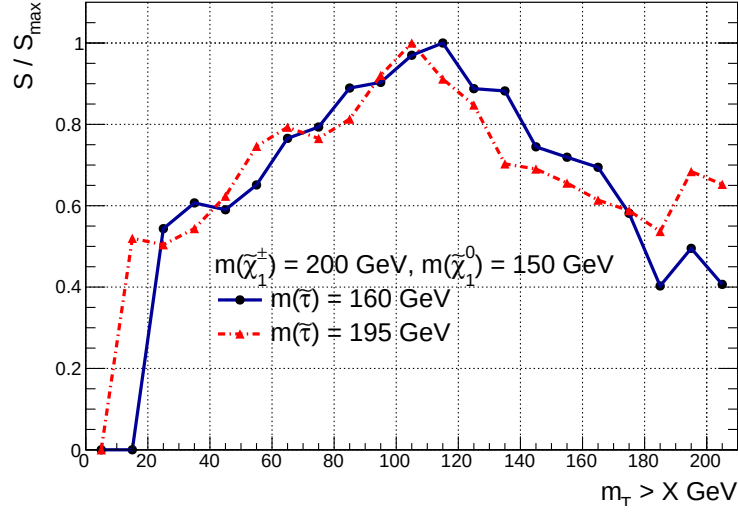


Figure 6.10: Signal significance $S = \frac{s}{\sqrt{s+B}}$ divided by the maximum significance S_{max} , as a function of m_T cut values for two signal samples 1) $m(\tilde{\chi}_1^\pm) = 200$ GeV, $m(\tilde{\tau}) = 160$ GeV, $m(\tilde{\chi}_1^0) = 150$ GeV and 2) $m(\tilde{\chi}_1^\pm) = 200$ GeV, $m(\tilde{\tau}) = 195$ GeV, $m(\tilde{\chi}_1^0) = 150$ GeV. The best signal significance is achieved for $m_T > 110$ GeV.

TeV are selected to form dijet candidates. We then choose the pair of jets with the highest dijet invariant mass as the VBF pair of jets.

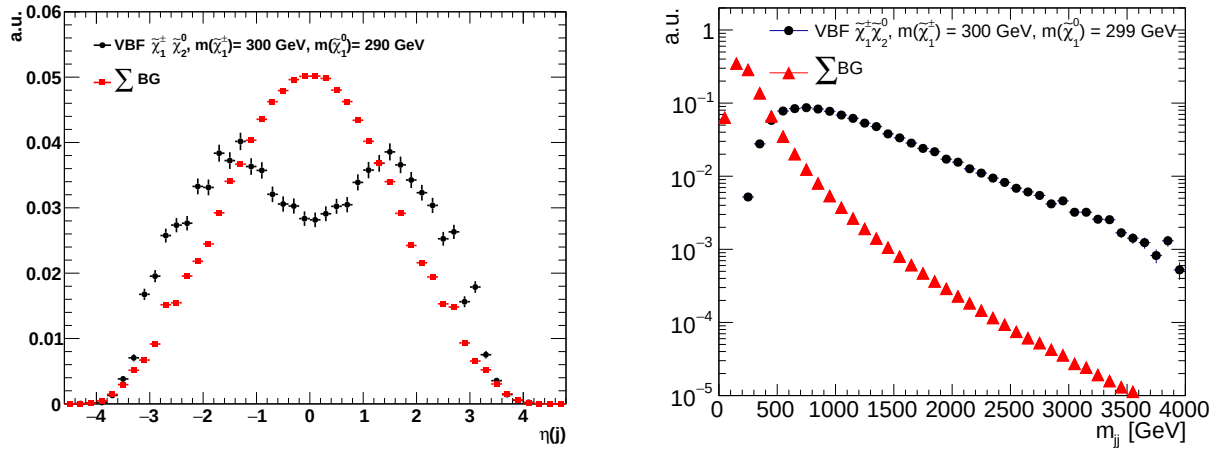


Figure 6.11: (Left) η distributions for jets normalized to unity for the three main backgrounds added together: W +jets, DY +jets and $t\bar{t}$ and one SUSY VBF signal sample $\tilde{\chi}_1^\pm \tilde{\chi}_2^0$ with $m(\tilde{\chi}_1^\pm) = 300$ GeV and $m(\tilde{\chi}_1^0) = 290$ GeV. (Right) Dijet mass distributions normalized to unity for the sum of backgrounds and one SUSY VBF signal sample $\tilde{\chi}_1^\pm \tilde{\chi}_2^0$ with $m(\tilde{\chi}_1^\pm) = 300$ GeV and $m(\tilde{\chi}_1^0) = 299$ GeV.

6.5 $t\bar{t}$ Background Estimation

As stated previously, we are searching for a SUSY signal with an event topology consisting of 2 large p_T jets with high pseudorapidity gap and large m_{jj} , one low- p_T muon, high E_T^{miss} and no b-jets. Even though, the no b-jet requirement is intended to eliminate the $t\bar{t}$ background, there is a remaining fraction of $t\bar{t}$ events that can fulfill the final state condition. As shown in Section 5.3, the b-jet tagging efficiency is about 69% for b-jets with $p_T > 30$ GeV and it is applied in the acceptance region of $|\eta| < 2.4$. There are $t\bar{t}$ events that could have their 2 b-jets out of the region of acceptance, or 1 b-jet out of the region of acceptance and the second within the region of acceptance but not identified due to the b-tagging inefficiency, a third possibility is to have $t\bar{t}$ events with both b-jets within the region of acceptance but still mis-identified by the b-tagger (as indicated in Figure 6.12). These type of $t\bar{t}$ events could have a W boson decaying to a muon and neutrino and the second W boson decaying to a pair of quarks that could mimic the VBF dijet condition. The dileptonic decays of $t\bar{t}$ events (both W bosons decaying leptonically) with one lepton misidentified and jets from initial or final state radiation, or the hadronic $t\bar{t}$ decays (both W bosons decaying hadronically) with one fake lepton could also emulate the VBF SUSY signal topology we are chasing. To estimate the number

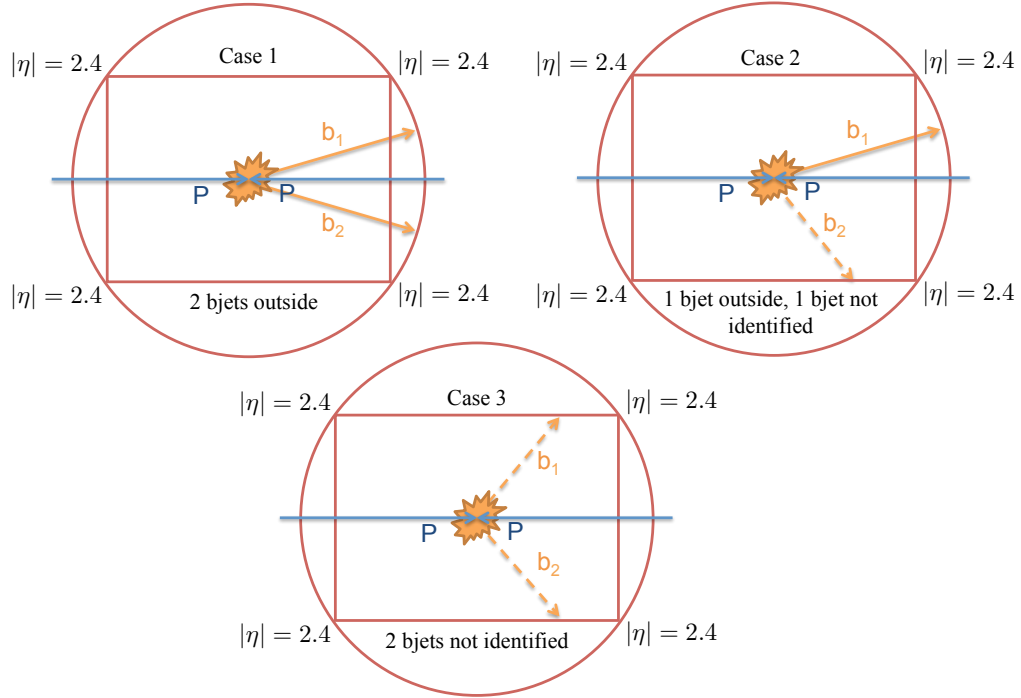


Figure 6.12: Scheme showing $t\bar{t}$ events with the production of 2 b-jets that are reconstructed as zero b-jets, the box inside the circle represents the detector acceptance region $|\eta| < 2.4$. There are three cases displayed: Case 1: events that could have their 2 b-jets out of the region of acceptance. Case 2: events with 1 b-jet out of the region of acceptance and the second within the region of acceptance but not identified. Case 3: events with both b-jets within the region of acceptance but still misidentified.

of $t\bar{t}$ background events, we use the strategy proposed in Section 6.1.1. We must identify a variable that allows us to define a data CR with enriched $t\bar{t}$ background. This choice should guarantee a high

purity of $t\bar{t}$ background compared to other backgrounds of the SM, and at the same time, it must ensure that the contamination due to the signal we search for is very low. In other words, when applying a requirement on this variable we should obtain in the CR a negligible number of events for the signal sample in comparison to the number of events obtained in the SR. For the $t\bar{t}$ case, the variable that meets all the previous requirements is the number of b-jets. This choice is made due to the b-jets produced in the decay of the top. The next step is to determine whether or not the choice of 1 b-jet as a variable to define our control regions has some bias in the kinematic distributions, in particular, those that are based on the selection of VBF (m_{jj} and $\Delta\eta_{jj}$). Figure 6.13 shows the comparison of the VBF m_{jj} and $\Delta\eta_{jj}$ distributions while Figure 6.14 shows the comparison of the muon kinematics (p_T and m_T distributions) in $t\bar{t}$ simulated events for central selections requiring 1 b-jet (CR1) and 0 b-jets. Given the good agreement between the distributions, we can conclude that the 1 b-jet requirement does not bias the lepton p_T or m_T distributions neither the jets m_{jj} and $\Delta\eta_{jj}$ distributions. In addition, we have to confirm that there is an equivalent composition

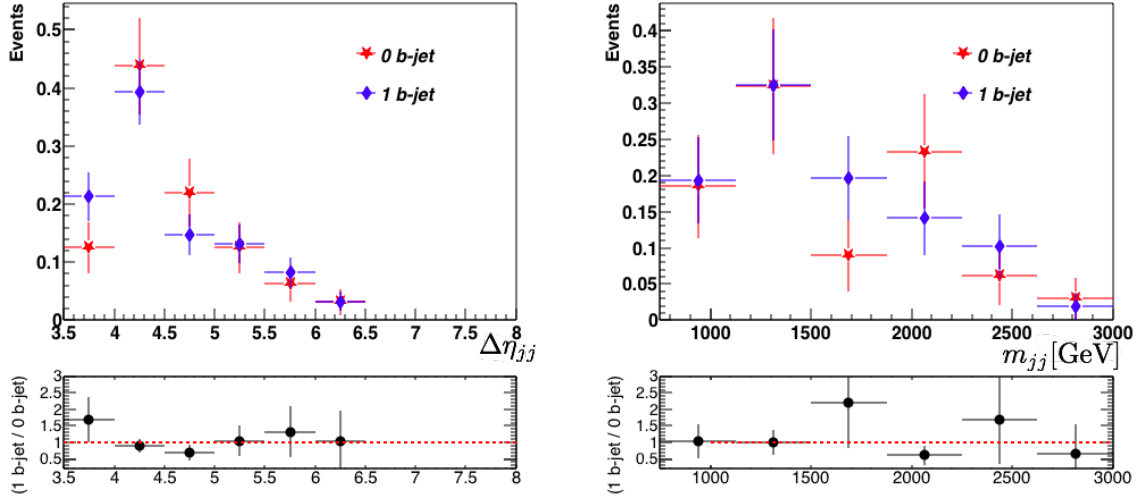


Figure 6.13: Comparison of the $t\bar{t}$ VBF $\Delta\eta_{jj}$ (left) and m_{jj} (right) distributions in events with 0 b-jets vs. 1 b-jet.

of real muons in the control and signal regions to guarantee that the scale factors are reflecting the proper corrections in the signal region. A muon produced in the MC generator that passes all muon selections is considered as a real muon, any other particle reconstructed as a muon but not identified at generation level as a muon is counted as a fake muon. We measure the fraction of real muons in the 0 b-jet and 1 b-jet regions and check if there is any difference. Table 6.4 provides these values, together with the VBF efficiency, determined as the fraction of events passing VBF requirements over total number of events before VBF requirements. The conclusion is that the muon content as well as the VBF efficiency are similar, within statistics, for both 0 b-jet and 1 b-jet samples. Note that 0 b-jet and 1 b-jet samples have contributions from b-jets produced out of the tracker acceptance: 1 b-jet with $|\eta| \geq 2.4$ and a second b jet with $|\eta| < 2.4$ (in one case untagged and in the other case tagged, respectively). A 2 b-jet sample could only have $t\bar{t}$ event topologies with the 2 b-jets being central and might produce some bias. To avoid any possible bias, we avoid requirements

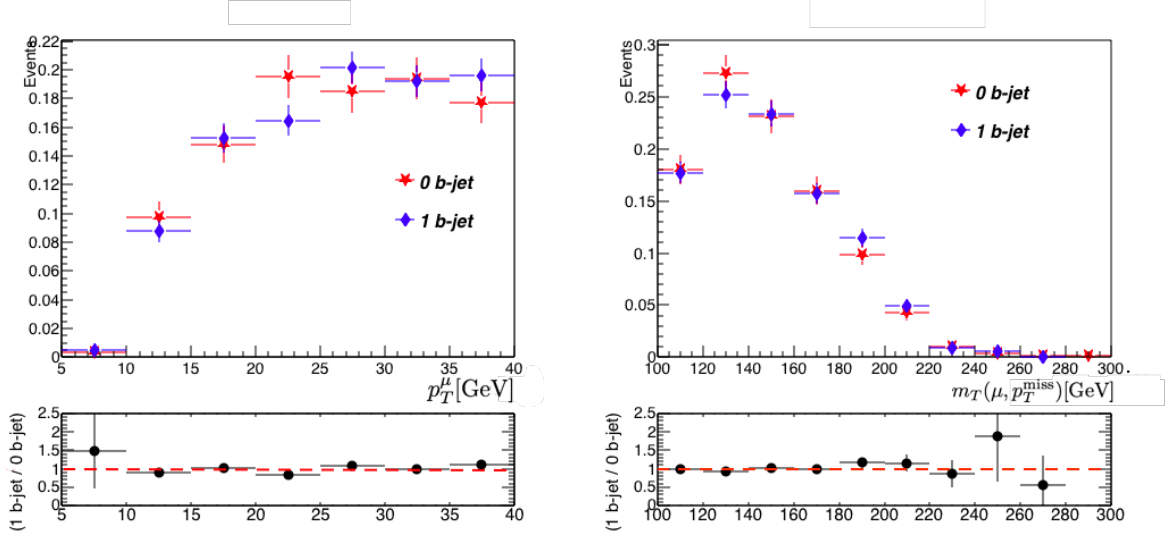


Figure 6.14: Comparison of the $t\bar{t}$ VBF p_T (left) and m_T (right) distributions in events with 0 b-jets vs. 1 b-jet.

Table 6.4: Comparison of the composition of $t\bar{t}$ events in the signal and control regions for real muons. MC generated muon that passes all identification requirements is counted as a real muon, any other MC generated particle that is reconstructed as a muon is counted as a fake.

Composition of muons	0 b-jets (signal region)	1 b-jets (control region)
Real Muons		
Total Muons = Real muons + Fakes	0.99 ± 0.01	0.98 ± 0.01
$\epsilon_{t\bar{t}}(\text{VBF cuts})$	$(6.05 \pm 1.15)\%$	$(6.62 \pm 0.66)\%$

that locate us farther away from the SR. In this case we just define the CR's by requiring exactly 1 b-jet. We have checked anyway, that the difference in statistics between the condition of exactly 1 b-jet and the condition of at least 1 b-jet is only 1.7% in the uncertainty of the SF. Figure 6.15 describes how the SR, CR1 and CR2 are defined.

Figure 6.16 shows the E_T^{miss} , m_T , $\eta(\mu)$, and $p_T(\mu)$ distributions for CR1 $t\bar{t}$ sample. The magenta bands represent the statistical uncertainties due to the sum of all the MC samples. Given the good agreement on data/MC ratio, we can conclude that the MC is correctly describing the data, in particular for the variables of interest p_T and m_T that are the variables with requirements applied in the central selection. The SF^{CR1} is calculated as follows:

$$SF^{\text{CR1}} = \frac{N^{\text{CR1}}(\text{Data}) - N^{\text{CR1}}(\text{other BGs})}{N^{\text{CR1}}(t\bar{t})}. \quad (6.7)$$

Table 6.5 lists the number of observed events in data as well as the expected MC contributions in CR1. The uncertainties are based on the statistics of the MC samples. The measured data-to-MC scale factor for CR1 is $SF^{\text{CR1}} = 0.97 \pm 0.06$. The distributions shown in Figure 6.16 have already been corrected with this scale factor.

To understand the behaviour of data and MC after applying the VBF selection, CR2 is defined by maintaining the VBF selections and requiring one b-jet (similar to CR1, except with events satisfying

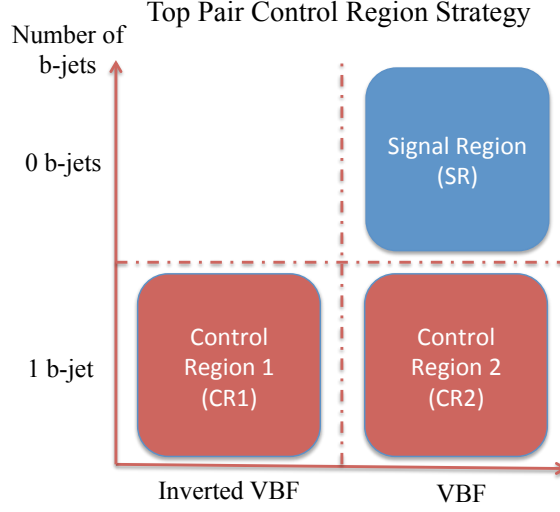


Figure 6.15: $t\bar{t}$ background estimation using CR1 and CR2.

Table 6.5: Predicted and observed background yields for control regions CR1 and CR2 defined to enhance $t\bar{t}$ background.

Sample	Yield CR1	Yield CR2
Diboson	3.7 ± 1.0	—
QCD	3.2 ± 2.3	—
Single Top	48.5 ± 2.7	2.0 ± 0.5
$W + \text{jets}$	33.1 ± 2.6	0.5 ± 0.1
$DY + \text{jets}$	1.2 ± 0.1	—
$t\bar{t}$	438.3 ± 11.4	28.1 ± 2.9
Total MC	528.0 ± 12.3	30.6 ± 3.0
Purity	83.0%	92.1%
Data	517	25
Scale Factor	0.97 ± 0.06	0.80 ± 0.20

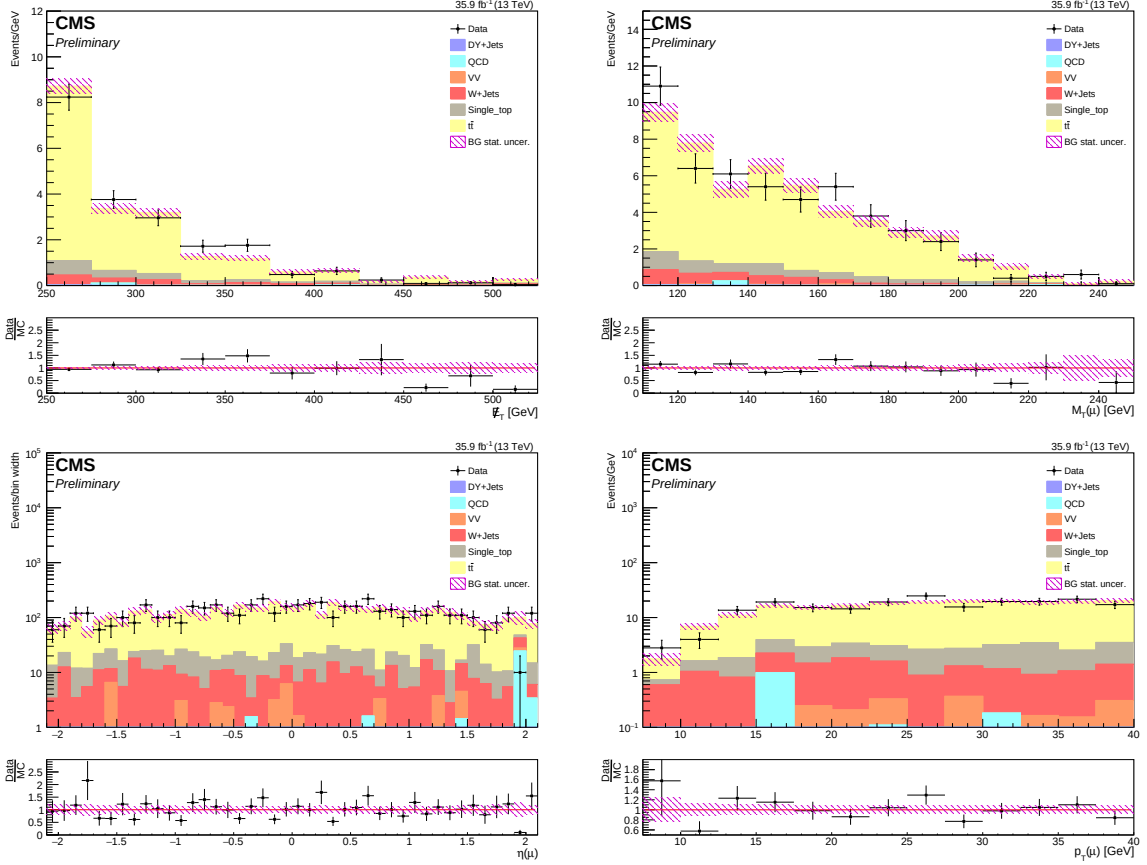


Figure 6.16: (a) E_T^{miss} , (b) $m_T(\mu, E_T^{\text{miss}})$, (c) $\eta(\mu)$, and (d) $p_T(\mu)$ distributions for $t\bar{t}$ CR1.

the VBF cuts). Table 6.5 shows the number of observed events in data as well as the expected MC contributions for CR2. Figure 6.17 shows the $\Delta\eta_{j,j}$ and m_{jj} distributions obtained for events in CR2. The scale factor for the VBF selection SF^{CR2} is calculated as follows:

$$SF^{\text{CR2}} = \frac{N^{\text{CR2}}(\text{Data}) - N^{\text{CR2}}(\text{other BGs})}{N^{\text{CR2}}(t\bar{t}) \cdot SF^{\text{CR1}}} \quad (6.8)$$

We obtain a value of 0.80 ± 0.20 . Note that the high uncertainty in this scale factor is due to the low statistics remaining after the VBF selection is applied. The number of expected $t\bar{t}$ events in the signal region is estimated by using the following equation:

$$N_{t\bar{t}}^{\text{Expected}} = N_{t\bar{t}}^{\text{MC}}(\text{SR cuts}) \cdot SF^{\text{CR1}} \cdot SF^{\text{CR2}} \quad (6.9)$$

where $N_{t\bar{t}}^{\text{MC}}(\text{SR cuts})$ is the predicted yield in MC simulation applying both the central selection and VBF requirement, SF^{CR1} is the correction factor for the central selection, obtained from the 1b + inverted VBF control samples, and SF^{CR2} the correction factor for the VBF requirement obtained from CR2. By using Equation 6.9 we get $N_{t\bar{t}}^{\text{Expected}} = 13.1 \pm 3.9$.

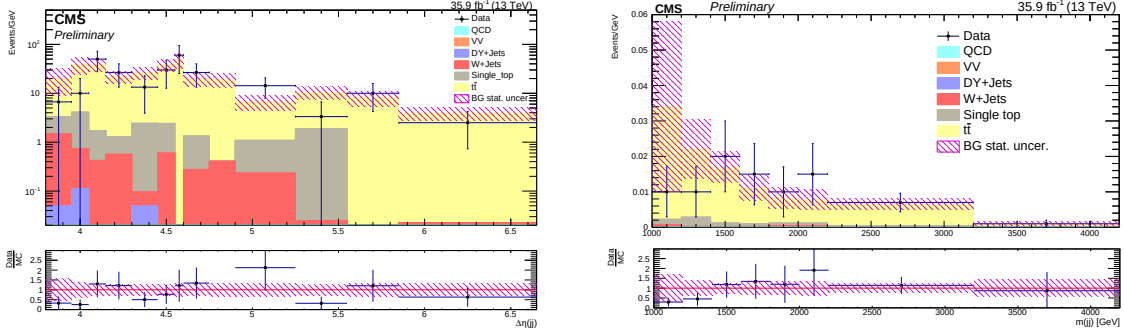


Figure 6.17: (a) $\Delta\eta(j, j)$, and (b) $m(j, j)$ distributions for $t\bar{t}$ CR2.

6.5.1 Validation Regions for $t\bar{t}$ background

Validation in a low- m_T region

Figure 6.18 shows validation regions at low m_T with the goal of making a closure test to the scale factors obtained in CR1 and CR2. One of the reasons for this validation is that the estimated value of $t\bar{t}$ events is obtained with a scale factor that involves data and MC (it is not completely data-driven). On the other hand, for the SF^{CR2} , where the VBF selection is applied, the statistics is low. These validation regions must contain a high purity of the studied background, and on the other hand, they must have higher statistics than the control regions. Additionally, they must be orthogonal to the control regions. The Validation Region 1 (VR1) is chosen by requiring m_T less than 110 GeV, inverted VBF, and exactly one b-jet. This region is used to validate the SF obtained in CR1. The way to validate the obtained result is to apply the SF^{CR1} as a correction to the MC of $t\bar{t}$ in VR1 and verify that there is good agreement between data and MC. A similar procedure is followed to validate the SF obtained in CR2. We define the Validation Region 2 (VR2) requiring m_T less than 110 GeV, passing VBF, and exactly one b-jet. Table 6.6 shows the number of events obtained in data and in MC for VR1 and VR2. According to simulation, $\sim 64\%$ of events in VR1 are $t\bar{t}$.

Figure 6.19 shows the $\eta(\mu)$ and $p_T(\mu)$ (with SF^{CR1} already applied) distributions for $t\bar{t}$ events belonging to region VR1. The distributions for the muon η and p_T variables for the VR1 show good agreement between data and corrected MC, with data-to-MC ratio of 1.01 ± 0.01 . A similar strategy is used for the validation of the SF^{CR2} . Figure 6.20 shows the η , p_T , m_T and m_{jj} distributions in the $t\bar{t}$ validation sample VR2 (after applying SF^{CR2}). Similar results as VR1 are obtained, the distributions show good agreement between data and MC.

$t\bar{t}$ background validation with Dimuon Samples

We carry out an additional test to demonstrate that the $t\bar{t}$ MC can be used to obtain the number of expected events in the signal region and correctly simulating the shapes of kinematic distributions in the signal region. We use a sample of dimuon events obtained from a muon dataset to which a single-muon trigger (HLT_IsoMu24_Eta2p1_v) is applied. Then we require $p_T(\mu) > 30$ GeV and a reconstructed b-jet. With this selection, we obtain an enriched sample of $t\bar{t}$ events and low signal contamination because the signal does not contain b-jets. Moreover, the topology of the signal does

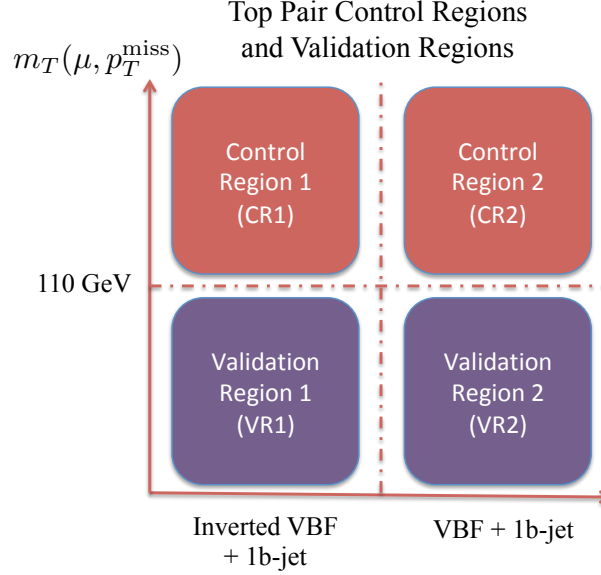
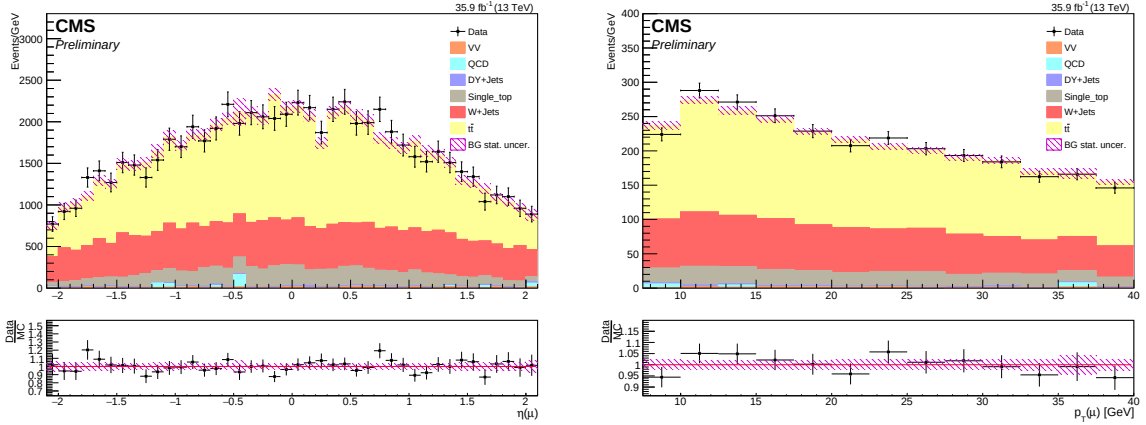


Figure 6.18: $t\bar{t}$ background Control Regions CR1 and CR2 with Validation Regions VR1 and VR2. The validation regions require $m_T < 110$ GeV, while the control regions require $m_T > 110$ GeV.



not produce two muons that can pass a p_T threshold of 30 GeV. Figure 6.21 shows the distributions of m_{jj} and m_T for the $t\bar{t}$ MC and the data with the previous selections. The ratio plots displayed in the lower part of Figure 6.21 show a good agreement between data and MC. This result confirms that we can use the $t\bar{t}$ MC with confidence to estimate the number of expected events in the signal region.

6.6 W+jets Background Estimation

Figure 6.22 describes the W+jets background control region strategy. W+jets events can mimic the

Table 6.6: Predicted and observed rates for the $t\bar{t}$ validation regions VR1 (VR2) obtained with 1 b-jets, $m_T < 110$ GeV and inverted (passing) VBF selections. They are used to validate scale factors measured from $t\bar{t}$ control samples: CR1 and CR2.

Sample	Yields for $t\bar{t}$ VR1	Yields for $t\bar{t}$ VR2
Diboson	37.4 ± 3.1	0.5 ± 0.4
QCD	38.4 ± 17.2	3.0 ± 2.3
Single Top	667.1 ± 8.9	29.0 ± 1.6
$W + \text{jets}$	2037.5 ± 22.9	60.8 ± 5.6
DY + jets	51.1 ± 2.0	2.5 ± 0.8
$t\bar{t}$	4457.0 ± 36.5	168.0 ± 7.1
Total MC	7288.5 ± 47.4	263.8 ± 9.5
Purity	61.2%	63.7%
Data	6861	288

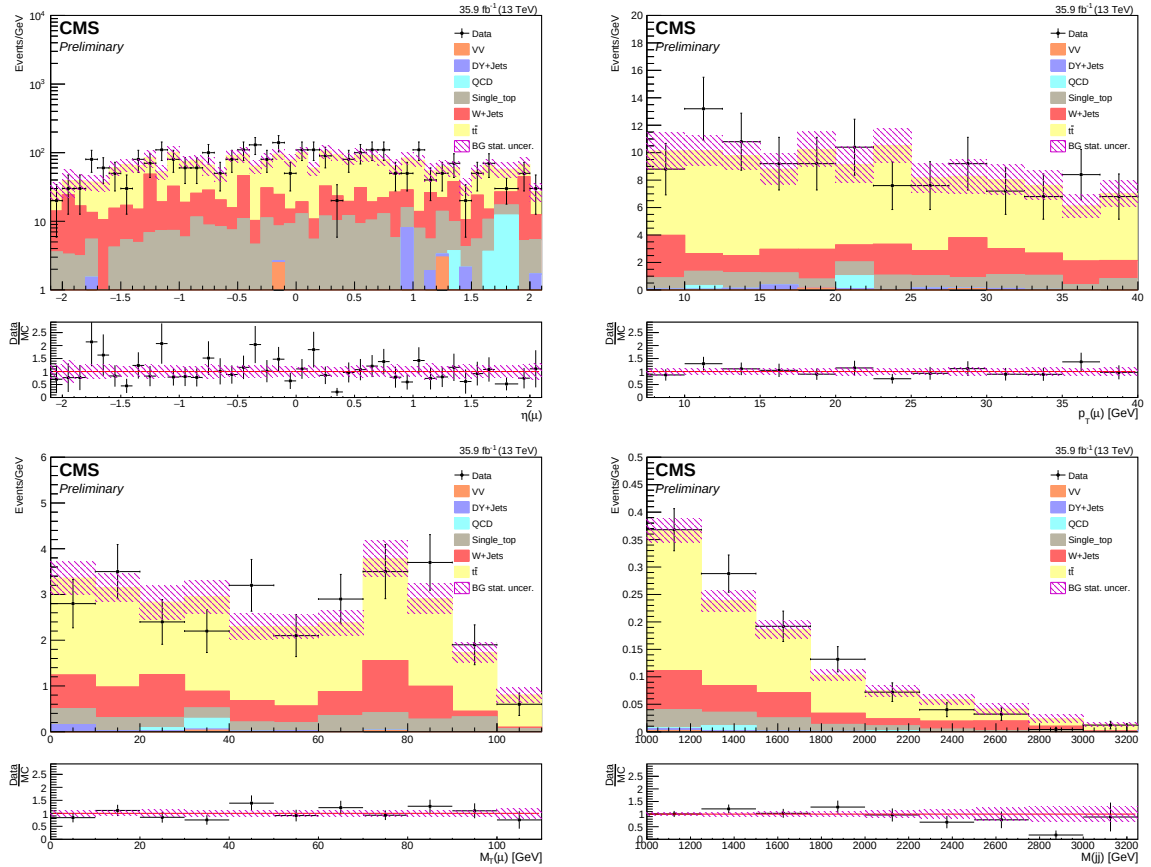


Figure 6.20: (a) $\eta(\mu)$, (b) $p_T(\mu)$, (c) $m_T(\mu, E_T^{miss})$, and (d) m_{jj} distributions for the $t\bar{t}$ validation region VR2.

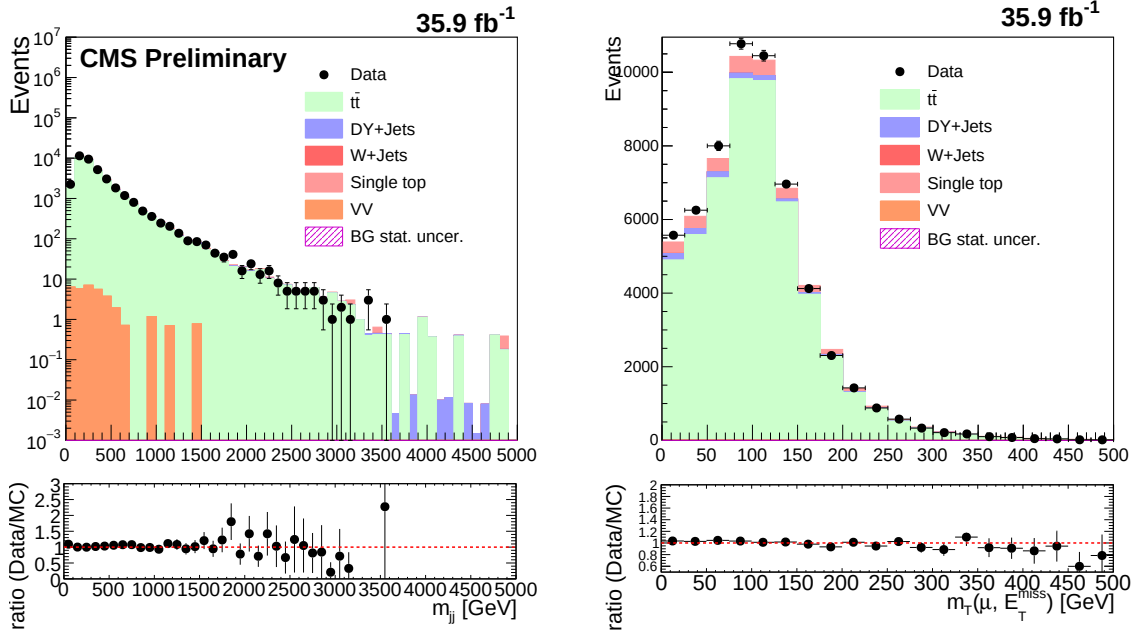


Figure 6.21: m_{jj} (left) and m_T (right) distributions for dimuon events in Data: obtained from a muon dataset (single-muon trigger applied), and requiring $p_T(\mu) > 30$ GeV and exactly one reconstructed b-jet. And MC: all the previous conditions with the exception of the trigger requirement. The purpose of all the selections is to obtain an enriched sample of $t\bar{t}$. The ratio plots show a good agreement between data and MC.

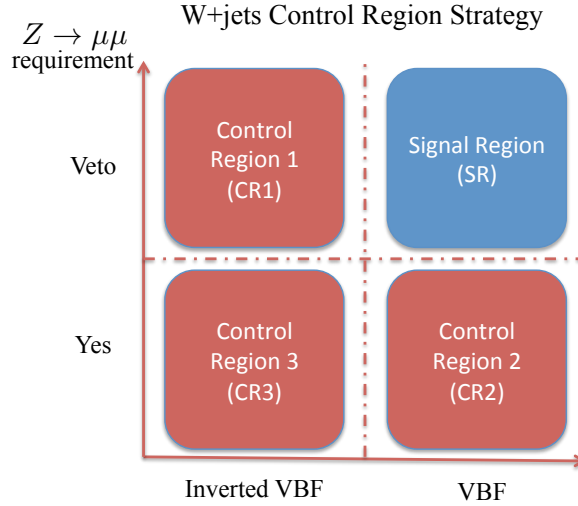


Figure 6.22: W+jets control region strategy.

SUSY signal topology we are investigating, because the leptonic decay of the W could produce a soft muon and MET and the associated jets could mimic the VBF topology. We follow a similar strategy

as the one used for $t\bar{t}$ background estimation, except that the W+jets events do not have b-jets. The challenge is to find orthogonal regions to measure the SF for the central selections (SF^{CR1}) and the SF for the VBF selection (SF^{CR2}). To obtain CR1, we apply all the central selections and inverted VBF. With these requirements, we obtain an enriched sample of W+jets with a purity of about 57%. In this region we calculate the SF^{CR1} in the usual way:

$$SF^{CR1} = \frac{N^{CR1}(\text{Data}) - N^{CR1}(\text{other BGs})}{N^{CR1}(\text{W+jets})}. \quad (6.10)$$

To obtain the SF for the VBF selection, we use the fact that the masses of the Z and W bosons are relatively close. So we pick two muons to reconstruct the Z mass by using the following requirements: two opposite sign muons with $p_T > 30$ GeV, a dimuon mass between 60 and 120 GeV, and all the central cuts. With these selections, we obtain an enriched sample of $Z \rightarrow \mu\mu$. Now we define two new control regions (CR2 and CR3 as shown in Figure 6.22). In CR3 we apply inverted VBF selections, and in CR2 we apply VBF selections. From CR3 we obtain a SF using the following equation:

$$SF^{CR3} = \frac{N^{CR3}(\text{Data}) - N^{CR3}(\text{other BGs})}{N^{CR3}(Z \rightarrow \mu\mu)}. \quad (6.11)$$

Then in CR2, we calculate a SF in the following way:

$$SF^{CR2} = \frac{N^{CR2}(\text{Data}) - N^{CR2}(\text{other BGs})}{N^{CR2}(Z \rightarrow \mu\mu) \cdot SF^{CR3}} \quad (6.12)$$

In the previous equation we are using the SF obtained in CR3 to correct the MC of $Z \rightarrow \mu\mu$ in CR2. With this, we have a SF that serves to correct the mismodeling produced by the VBF selection (SF^{CR2}). To calculate the number of expected events in the SR for the MC of W+jets we use the following equation:

$$N_{W+jets}^{\text{Expected}} = N_{W+jets}^{\text{MC}}(\text{SR cuts}) \cdot SF^{CR1} \cdot SF^{CR2} \quad (6.13)$$

Table 6.7 shows all the yields obtained for each control region. Figures 6.23 , 6.24 and 6.25 show the distributions for m_T and $p_T(\mu)$ for CR1, $m_{\mu\mu}$ and m_{jj} for CR2, and $m_{\mu\mu}$ and $p_T(\mu)$ for CR3 respectively. The excellent agreement observed in all the figures show that the SF's can be used to calculate the number of expected events with confidence. According to equation 6.13 the number of expected events is $N_{W+jets}^{\text{Expected}} = 14.7 \pm 3.4$.

6.6.1 Validation Regions for W+jets

We perform a validation process similar to that performed with $t\bar{t}$, using two regions of low- m_T as shown in Figure 6.26. To carry out the validation, we choose regions that are entirely orthogonal and have a high purity of W + jets events. The Validation Region 1 (VR1) is chosen by requiring $m_T < 110$ GeV and inverted VBF. This region is used to validate the SF obtained in CR1. The way to validate the SF is to apply the SF^{CR1} as a correction to the MC of W+jets in VR1 and verify that there is good agreement between data and MC. A similar procedure is followed to validate the SF obtained in CR2. We define the Validation Region 2 (VR2) requiring $m_T < 110$ GeV and passing VBF. Table 6.8 shows the number of events obtained in data and MC for VR1 and VR2. Figures 6.27 and 6.28 show the distributions $\eta(\mu)$ and $p_T(\mu)$ for VR1, and $\eta(\mu)$, $p_T(\mu)$, E_T^{miss} and m_{jj} for VR2 respectively. The excellent agreement observed in all the figures show that the SF's obtained in CR1 and CR2 can be used with confidence to correct the W+jets background yield for any mismodeling of the central selections and VBF selections.

Table 6.7: Predicted and observed events for the W+jets control region with inverted VBF selections CR1, for the $Z(\rightarrow \mu\mu) + \text{jets}$ control regions CR2 and CR3. CR1 is used to obtain a scale factor for central selections, SF^{CR1} . CR2 is used to obtain a scale factor for the VBF selection, SF^{CR2} . CR3 is used to obtain an scale factor for the central selections (SF^{CR3}) to correct the MC of $Z \rightarrow \mu\mu$ in CR2.

Sample	W+jets CR1 $_{\mu}$	$Z(\rightarrow \mu\mu) + \text{jets}$ CR2	$Z(\rightarrow \mu\mu) + \text{jets}$ CR3
Diboson	83.6 ± 5.3	25.0 ± 2.8	18533.7 ± 78.1
QCD	1.6 ± 0.7	373.3 ± 250.5	62412.8 ± 23860.0
Single Top	38.6 ± 2.7	20.3 ± 1.9	5625.1 ± 32.9
W + jets	619.3 ± 42.4	3.0 ± 0.9	1372.1 ± 100.9
DY + jets	39.2 ± 10.4	3100.7 ± 42.6	15795222.6 ± 7659.0
$t\bar{t}$	310.0 ± 11.1	252.5 ± 9.9	34336.6 ± 115.6
Total MC	1091.5 ± 45.4	3774.8 ± 254.3	15917502.9 ± 25059.7
Purity	56.7%	82.1%	99.2%
Data	1160	4174	15715632
Scale Factor	1.11 ± 0.10	1.13 ± 0.08	0.95 ± 0.001

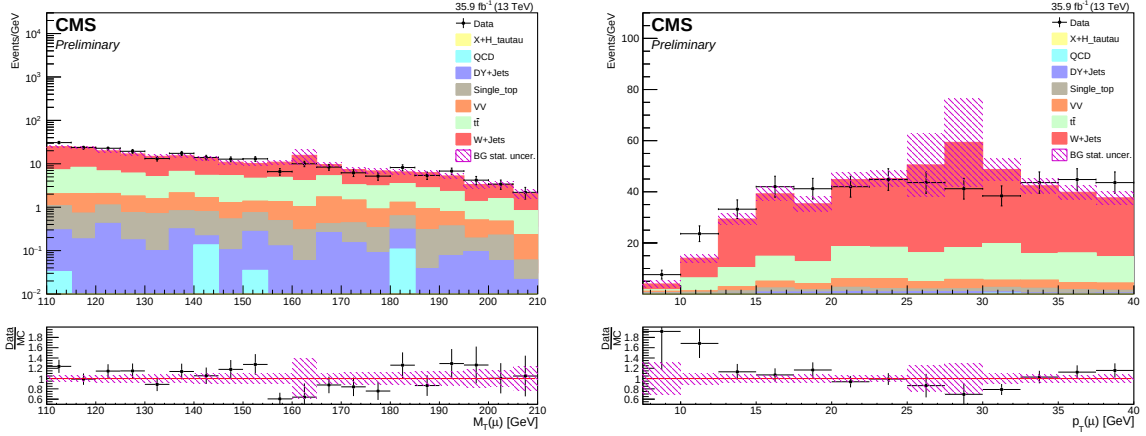


Figure 6.23: (a) $m_T(\mu, E_T^{\text{miss}})$ and (b) $p_T(\mu)$ distributions for W+jets CR1.

6.7 Other Backgrounds

Diboson production, DY+jets, Single Top and QCD may have final states similar to the topology of our signal. Figure 6.29 and Figure 6.30 show leading order Feynman diagrams for Diboson and QCD production that contribute to our signal region. However, all these processes contribute only with 12% of the total background yield (as shown in Table 6.9). $t\bar{t}$ and W+jets contribute with the remaining 88%, and this is the reason why they have a dedicated background estimation strategy.

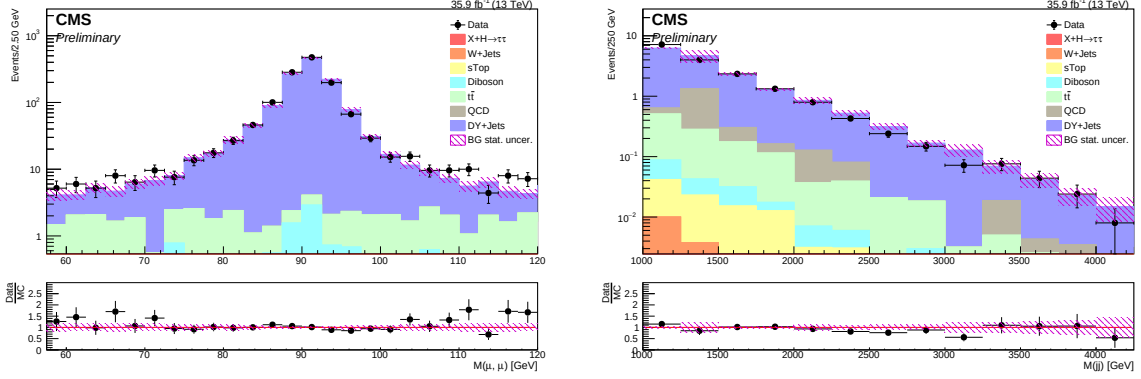


Figure 6.24: (a) $m_{\mu\mu}$ and (b) m_{jj} distributions for CR2.

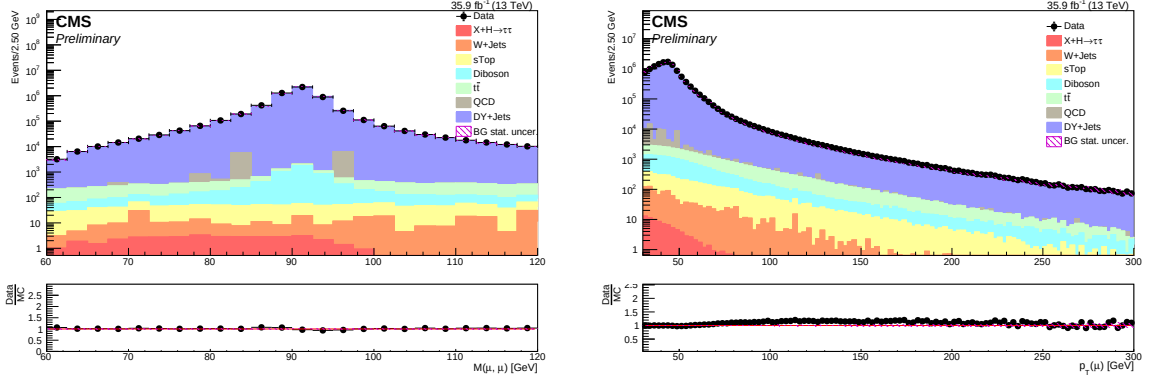


Figure 6.25: (a) $m_{\mu\mu}$ and (b) $p_T(\mu)$ distributions for Z + jets CR3.

6.8 Comparison of data and Background yields in the signal region

Table 6.9 shows the predicted and observed yields without any correction and with corrections applied to W+jets and $t\bar{t}$ background yields for the Signal region (SR). For the remaining backgrounds (different to W+jets and $t\bar{t}$) no corrections are applied to MC, given their low statistical contribution in the SR. Figure 6.31 shows the distribution of $\Delta\phi(\mu, p_T^{\text{miss}})$, p_T^{miss} , m_{jj} , jet multiplicity and $p_T(\mu)$ for data and MC in the SR.

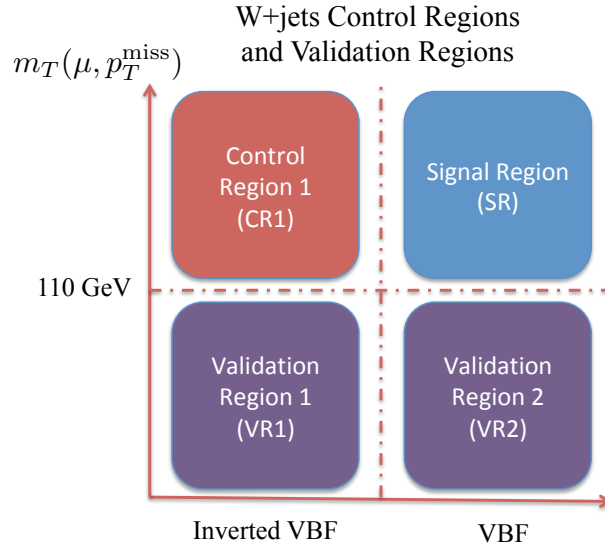


Figure 6.26: W+jets validation strategy.

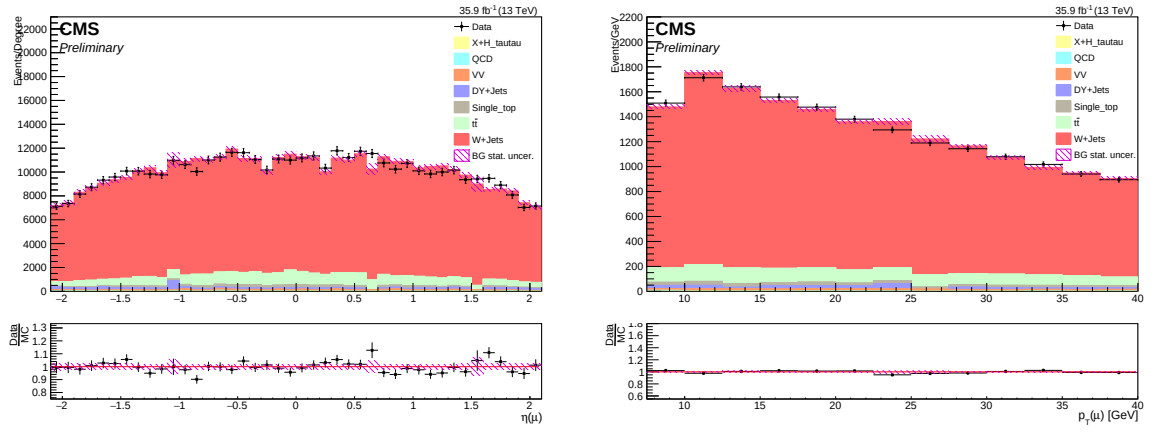


Figure 6.27: (a) $\eta(\mu)$ and (b) $p_T(\mu)$ distributions for the W+jets validation region VR1 $_{\mu}$.

Table 6.8: Predicted and observed yields for the W+jets validation region VR1 (VR2) obtained with $m_T < 110$ GeV and inverted (passing) VBF selections.

Sample	W+jets VR1 $_{\mu}$	W+jets VR2 $_{\mu}$
Diboson	516.6 ± 13.4	8.1 ± 1.7
QCD	56.6 ± 14.3	—
Single Top	730.0 ± 10.4	25.4 ± 1.8
W + jets	36550.6 ± 2260.9	777.9 ± 21.7
DY + jets	627.5 ± 43.8	14.9 ± 1.8
$t\bar{t}$	3325.3 ± 36.2	124.5 ± 7.0
Total MC	41820.2 ± 2261.7	950.8 ± 23.0
Purity	87.4%	81.8%
Data	42088.0	969
Scale Factor	1.01 ± 0.06	—

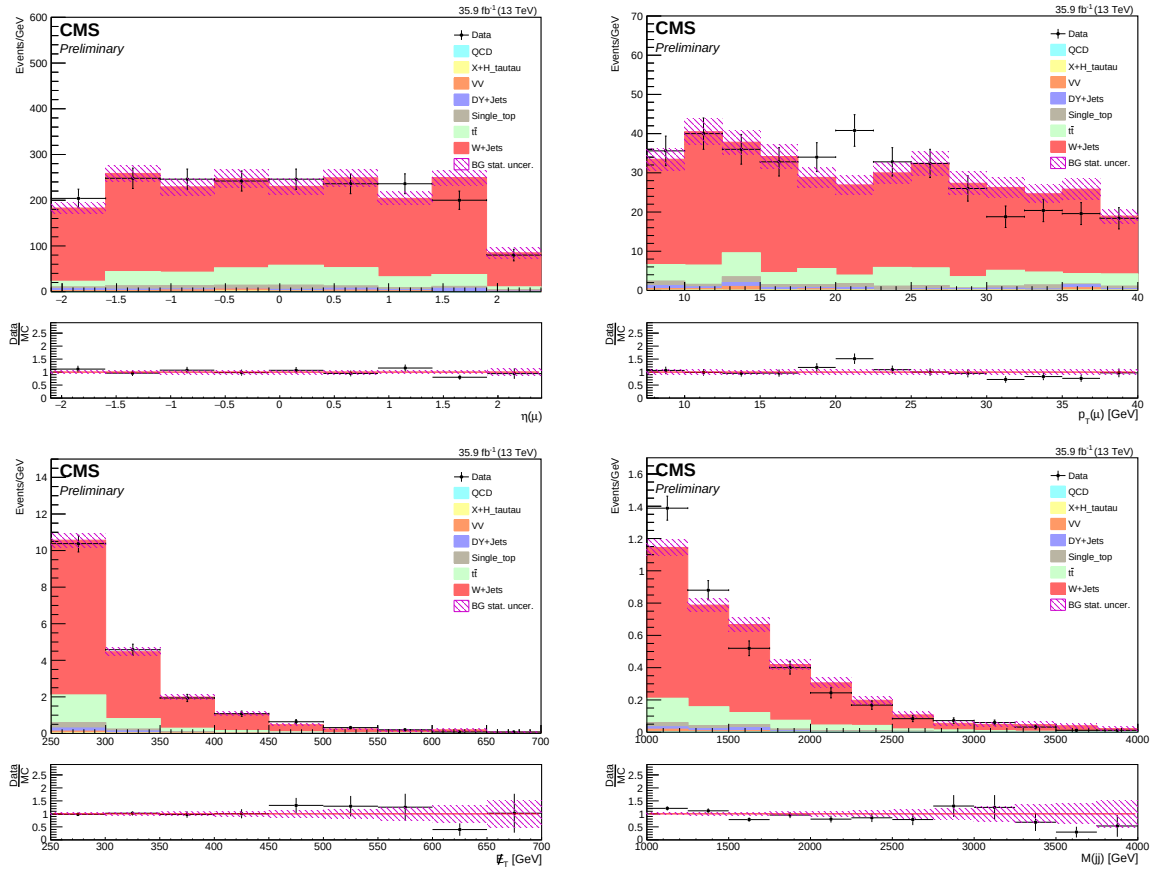


Figure 6.28: (a) $\eta(\mu)$, (b) $p_T(\mu)$, (c) E_T^{miss} , and (d) m_{jj} distributions for the W+jets validation region VR2 $_{\mu}$.

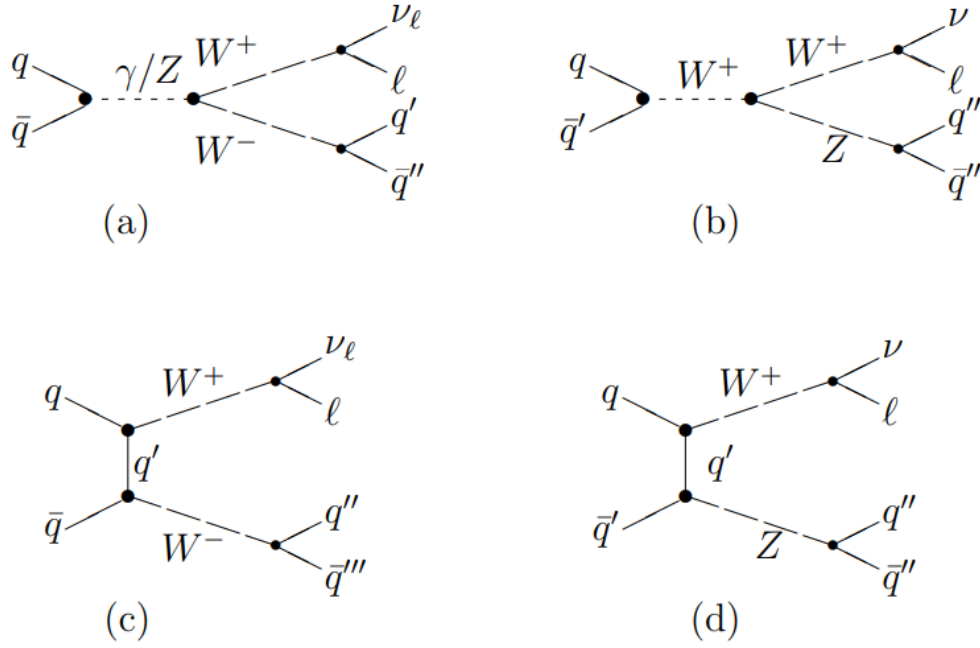


Figure 6.29: Leading-order Feynman diagrams for Diboson production in proton-proton collisions that have final states similar to our VBF signal topology.

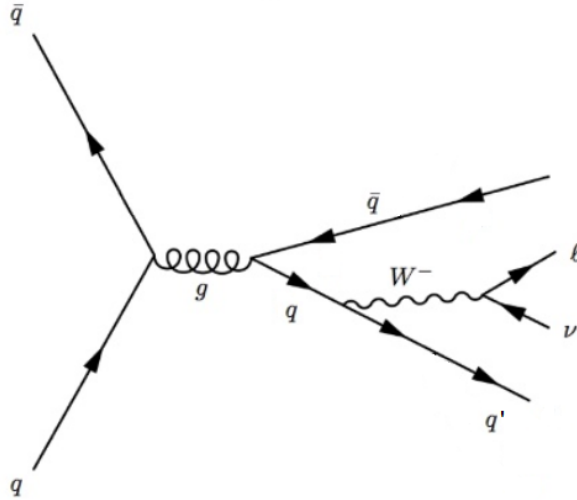


Figure 6.30: Leading-order Feynman diagram for QCD proton-proton collisions that has final states similar to our VBF signal topology. The diagram displayed is just one example of how a QCD process can generate the same final state we are studying.

Table 6.9: Predicted and observed yields in the signal region without any correction and with corrections applied to W+jets and ttbar backgrounds yield. As an example, the $t\bar{t}$ background (BG) prediction is calculated as follows: $N_{t\bar{t}}^{\text{Corrected}} = N_{t\bar{t}}^{\text{MC}}(\text{SR cuts}) \cdot SF^{\text{CR1}} \cdot SF^{\text{CR2}}$

Sample	$N_{\text{BG}}^{\text{MC}}(\text{SR})$	$N_{\text{BG}}^{\text{Corrected}}$
Diboson	1.3 ± 0.6	1.3 ± 0.6
QCD	—	—
Single Top	2.3 ± 0.7	2.3 ± 0.7
$W + \text{jets}$	10.2 ± 2.1	13.2 ± 3.1
$DY + \text{jets}$	0.2 ± 0.1	0.2 ± 0.1
$t\bar{t}$	16.9 ± 2.6	13.1 ± 3.9
Total BG Prediction	30.9 ± 3.5	31.6 ± 5.3
Data	—	36

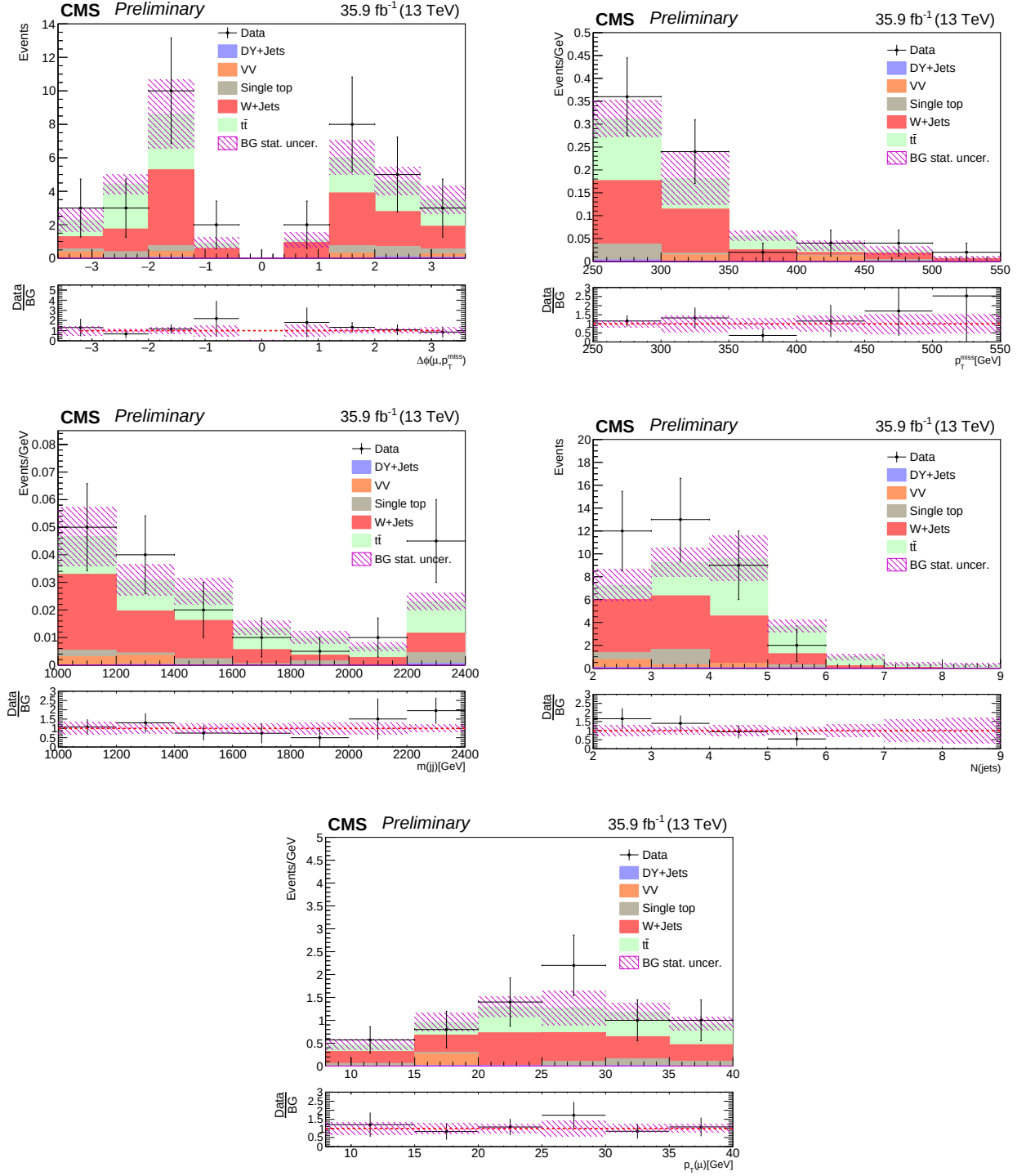


Figure 6.31: (top left) $\Delta\phi(\mu, p_T^{\text{miss}})$, (top right) p_T^{miss} , (middle left) m_{jj} , (middle right) jet multiplicity, and (bottom left) $p_T(\mu)$ distributions in the SR. Only the W+jets and ttbar background contributions have scale factors applied.

Chapter 7

Results

Motivated by chargino/neutralino production through vector boson fusion within R-parity conserving MSSM, for scenarios where $\tilde{\chi}_1^\pm/\tilde{\chi}_2^0$ could decay to lighter sleptons, and where the mass splitting between chargino/neutralino and LSP is small (compressed spectra), we embarked ourselves in the mission of searching for an excess of events in a signal region defined by one single muon, E_T^{miss} and 2 VBF jets in the final state (VBF jj + μ + E_T^{miss}). The VBF jets are required to be a pair of jets with very high invariant mass ($m_{jj} > 1$ TeV) and separated by a large pseudorapidity gap ($|\Delta\eta_{jj}| > 3.8$). We estimate, by MC simulations, the backgrounds that survive the set of requirements (central selections + VBF) motivated by the enhancement of a hypothetical signal for the VBF SUSY scenario under consideration. The main sources of irreducible backgrounds in the SR are: $t\bar{t}$ and W+jets events. We first verify that the MC describes correctly the different data kinematic distributions in control regions (enriched by each background type), orthogonal to the SR. We then apply two scale factors to the number of MC events in the signal region, to correct any MC mismodeling for the applied requirements. SF^{CR1} for the central selections and SF^{CR2} for the VBF selection. Table 7.1 shows the contributions of the SM backgrounds that have an impact on the SR, and provides the event yields before and after applying the scale factors to $t\bar{t}$ and W+jets MC events. As it can be observed from Table 7.1, the contributions from other backgrounds (such as single top, diboson, DY+jets) are very small. Additionally, six SUSY VBF signal samples with different values for $m(\tilde{\chi}_1^\pm)$, $m(\tilde{l})$ and $m(\tilde{\chi}_1^0)$ are presented. When the mass difference between $m(\tilde{\chi}_1^\pm)$ and $m(\tilde{\chi}_1^0)$ is 10, 30 and 50 GeV. Figure 7.1 shows m_T and m_{jj} distributions for data and MC. No data excess is observed over the expected background yields. For comparison, the distributions from two SUSY VBF signal samples are also indicated: $m(\tilde{\chi}_1^\pm) = 200$ GeV, $m(\tilde{l}) = 185$ GeV, $m(\tilde{\chi}_1^0) = 170$ GeV and $m(\tilde{\chi}_1^\pm) = 300$ GeV, $m(\tilde{l}) = 285$ GeV, $m(\tilde{\chi}_1^0) = 270$ GeV. In the following sections we compare the outcome of our analysis in the muon channel to the results from electron, tau and invisible channel and then provide exclusion limits.

7.1 Comparison with electron and tau channel

The same methodology as explained for the muon channel is applied for the electron and tau channels. The final state topology is VBF jj + ℓ + MET, where ℓ is a single isolated soft p_T lepton (either electron, muon or tau). All selections applied to each leptonic channel are summarized in Table 7.2. Table 7.3 shows, for the electron and tau channels, relevant values for the SR: the raw MC yields,

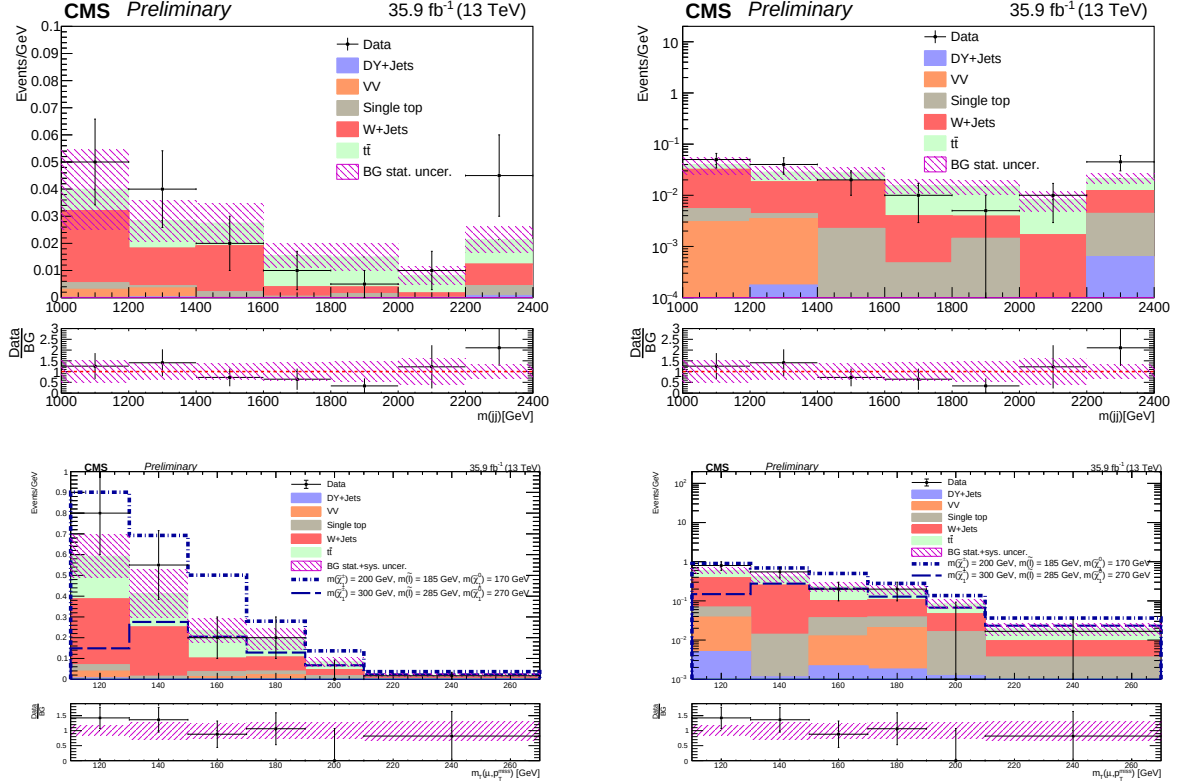


Figure 7.1: Normal scale (in the left) and log vertical scale (in the right) for m_{jj} and m_T distributions in the signal region. For comparison two signal samples are presented: 1) $m(\tilde{\chi}_1^\pm) = 200$ GeV, $m(\tilde{l}) = 185$ GeV, $m(\tilde{\chi}_1^0) = 170$ GeV and 2) $m(\tilde{\chi}_1^\pm) = 300$ GeV, $m(\tilde{l}) = 285$ GeV, $m(\tilde{\chi}_1^0) = 270$ GeV. In all plots shown, the last bin is the overflow bin.

the scale factors to correct MC mismodeling of central selections (SF^{CR1}) and MC mismodeling of VBF requirements (SF^{CR2}), and the final expected background yields. Note, that scale factors are only applied on the main sources that account for about 90% of the total background yield in the SR. Figures 7.2 and 7.3 show the m_T distributions for the electron and tau channel, respectively. A radical difference in the tau channel, compared to electron and muon, is the QCD background. While for the muon and electron channel the contribution of QCD is negligible, in the tau channel QCD is the dominant source. This difference is due to a greater probability of obtaining tau fakes in the reconstruction process, since the tau is reconstructed from its hadronic decay. No event excess is observed in neither muon, electron nor tau channels. Given the amount of data analyzed, we then establish 95% CL exclusion limits on the SUSY VBF signal we scrutinized.

7.2 Invisible Channel

An additional search considered in a complementary analysis is the case where none of the produced leptons in the final state can be reconstructed. The topology of the final state, in this situation, is MET accompanied by the VBF jets. This type of signals would be expected in very compressed

Table 7.1: MC predicted background contributions in the signal region for the VBF $jj + \mu + E_T^{\text{miss}}$ channel. As an example, the $t\bar{t}$ background (BG) prediction is calculated as follows: $N_{t\bar{t}}^{\text{Data}} = N_{t\bar{t}}^{\text{MC}}(\text{SR cuts}) \cdot SF^{\text{CR1}} \cdot SF^{\text{CR2}}$.

Sample	$N_{\text{BG}}^{\text{MC/Data}}(\text{SR})$	SF^{CR1}	SF^{CR2}	$N_{\text{BG}}^{\text{Data}}$
Diboson	1.3 ± 0.6	—	—	1.3 ± 0.6
QCD	—	—	—	—
Single Top	2.3 ± 0.7	—	—	2.3 ± 0.7
$W + \text{jets}$	10.2 ± 2.1	1.10 ± 0.10	1.18 ± 0.09	13.2 ± 3.1
DY + jets	0.2 ± 0.1	—	—	0.2 ± 0.1
$t\bar{t}$	16.9 ± 2.6	0.97 ± 0.06	0.80 ± 0.20	13.1 ± 3.9
BG Prediction	30.9 ± 3.5	—	—	31.6 ± 5.3
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 200, 195, 190 \text{ GeV}$	—	—	—	26.2 ± 0.6
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 200, 185, 170 \text{ GeV}$	—	—	—	52.4 ± 0.8
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 200, 175, 150 \text{ GeV}$	—	—	—	47.2 ± 0.7
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 300, 295, 290 \text{ GeV}$	—	—	—	6.6 ± 0.1
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 300, 285, 270 \text{ GeV}$	—	—	—	17.9 ± 0.2
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 300, 275, 250 \text{ GeV}$	—	—	—	17.0 ± 0.2
Data	—	—	—	36

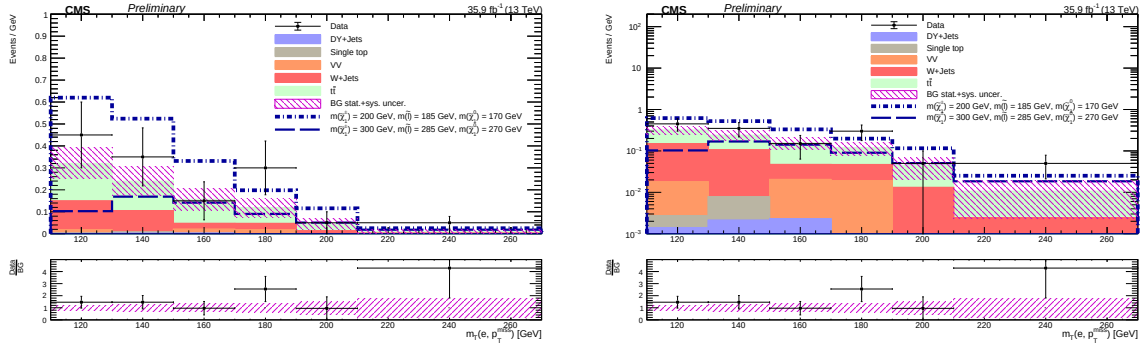


Figure 7.2: m_T distribution in the VBF $jj + e + E_T^{\text{miss}}$ signal region (left is normal scale, while right is log vertical scale). The band in the ratio plot represents the total uncertainty on the background prediction (systematic and statistical).

spectra supersymmetric scenarios, where the mass differences in the chargino/neutralino system are very small ($\Delta m = m_{\tilde{\chi}_1^\pm} - m_{\tilde{\chi}_1^0} < 10 \text{ GeV}$) implying the leptons of the final state would not have enough energy to be reconstructed by the detector. In this case, unlike searches with leptons, the variable to explore this signal topology is m_{jj} , since there is no reconstructed lepton available to do an m_T shape analysis. Table 7.4 depicts the selection of cuts used in the invisible channel while Figure 7.4 shows the distribution of m_{jj} for data and MC backgrounds in the SR. As in the cases with a reconstructed lepton in the final state, no excess of events is observed.

		$e^\pm$ Channel	$\mu^\pm$ Channel	$\tau^\pm$ Channel
lepton flavor	$e^\pm$	Selected	Vetoed	Vetoed
	$\mu^\pm$	Vetoed	Selected	Vetoed
	$\tau^\pm$	Vetoed	Vetoed	Selected
lepton ID		Medium ID	Tight ID	Discrimination prong type 1 hps
lepton Iso		-	0.25	Tight Isolation MVA run2 v1 DB new DMwLT
lepton p_T	$e^\pm$	10-40 GeV	> 10 GeV	> 10 GeV
	$\mu^\pm$	> 8 GeV	8-40 GeV	> 8 GeV
	$\tau^\pm$	> 20 GeV	> 20 GeV	20-40 GeV
$m_T(l, p_T^{miss})$	$e^\pm$	> 110 GeV	No Cut	No Cut
	$\mu^\pm$	No Cut	> 110 GeV	No Cut
	$\tau^\pm$	No Cut	No Cut	> 110 GeV
p_T^{miss}		> 250 GeV		
Overlap Removal		Yes		
b-jet		$N_{bjet} = 0$, CSVv2M (Medium WP), $p_T > 30$ GeV, $ \eta < 2.4$		
VBF Selection		$N_{jets} \geq 2$ with $p_T^{jet} > 60$ GeV, $ \eta(j) < 5$ At least one set of jets (j_1, j_2) with $\eta(j_1) \times \eta(j_2) < 0$, $\Delta\eta(j_1, j_2) > 3.8$, $m_{jj} > 1$ TeV		
SR m_T Bins		[110-130], [130-150], [150-170], [170-190], [190-210], [210- ∞]		

Table 7.2: Summary of signal region selections for the three leptonic channels.

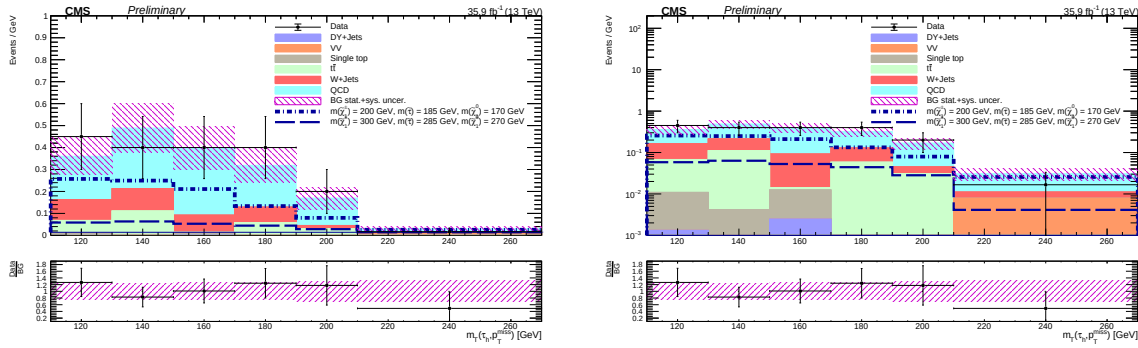


Figure 7.3: m_T distribution in the VBF $jj + \tau_h + E_T^{miss}$ signal region (left is normal scale, while right is log vertical scale). The band in the ratio plot represents the total uncertainty on the background prediction (systematic and statistical).

7.3 Systematic Errors

The following naming convention is used for Table 7.5 showing the systematics uncertainties considered in this analysis and their percentage contribution.

- **lumi:** uncertainty on the measured luminosity [214]
- **μ ID:** uncertainty in muon reconstruction and identification [215].
- **e ID:** uncertainty in electron reconstruction and identification [216].
- **τ_h ID:** uncertainty in hadronic tau reconstruction and identification [217].

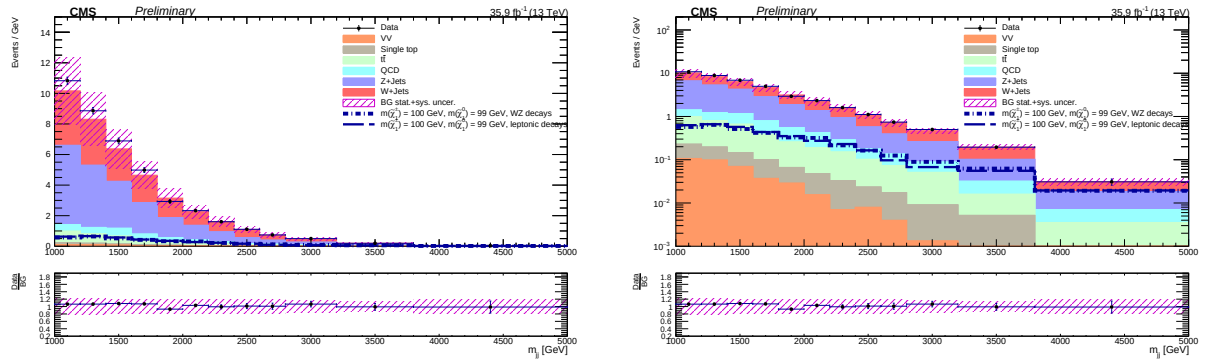


Figure 7.4: SR background and data distributions for dijet invariant mass in linear scale (left) and logarithmic scale (right) for the invisible channel (after applying the selection requirements from Table 7.4).

- **Trigger** : trigger efficiency uncertainty
- **b ID** : uncertainty in b-tagging reconstruction and identification [218, 219]
- **JES**: uncertainty in the calibration of the energy of jets in ECAL and HCAL
- **TES**: uncertainty due to τ_h energy scale [217].
- **MMS**: uncertainty due to muon momentum scale [215].
- **EES**: uncertainty due to electron energy scale [216].
- **pdf**: uncertainty due to variations in the parton distribution functions [220].
- **scale**: generator scale uncertainty
- **bin-by-bin stat.** : bin-by-bin shape uncertainty due to closure/validation tests
- **Closure+Norm.** uncertainty on the background normalization due to the data-driven estimate

7.4 Limits

Any search of new physics can be interpreted as a hypothesis test, between a null hypothesis H_0 that assumes there is no new physics and an alternative hypothesis H_1 that assumes the contributions of new physics in the variables measured. To pretend that there is a discovery, one must reject the null hypothesis with a certain probability, called the p -value. For the case of the null hypothesis, the number of expected events (λ) is considered as the background, $\lambda = b$, while for the alternative hypothesis, which includes the contribution of new physics, the number of expected events is $\lambda = s + b$. Where s is the signal. However, in our results we have a histogram with several bins (N) so that we can consider each bin as an independent experiment, then each bin has its own probability distribution. Since the number of independent results is small (number of bins) and the cross

section for a signal event is very low, compared with the cross-section for background events, we assume that the probability of observed events follow a Poisson distribution. Subsequently, we can build a Likelihood function as the product of the Poisson distribution for all bins.

$$\mathcal{L} = \prod_i^N \mathcal{L}(\lambda_i, k_i) = \prod_i^N \frac{\lambda_i^{k_i} e^{-\lambda_i}}{k_i!} \quad (7.1)$$

Note that the muon channel is not the only channel used to search SUSY VBF. If we want to combine the results taking into account all the leptonic channels (muon, electron, tau and neutrinos). Then the combined likelihood function is given by the product of the likelihood functions for each leptonic channel:

$$\mathcal{L}_{\text{combined}} = \mathcal{L}_\mu \mathcal{L}_e \mathcal{L}_\tau \mathcal{L}_\nu \quad (7.2)$$

This combined function is used to establish limits in the cross-section of the signal. The systematic uncertainties are introduced in the yields as nuisance parameters (θ). A nuisance parameter is an amount used to determine how much the values obtained from signal and background are affected due to the different systematics taken into account. Then the likelihoods depend on these nuisance parameters. We use the CLs method, which makes a test between the alternative hypothesis H_1 and the null hypothesis H_0 through the likelihood ratio. This ratio is considered to provide the best discriminator between the two hypotheses [221] and it is given by the following expression:

$$LR = -2 \ln \left(\frac{\mathcal{L}_{s+b}(\theta')}{\mathcal{L}_b(\theta'')} \right) \quad (7.3)$$

When the LR discriminator is positive, it means that the observed events are more compatible with the null hypothesis H_0 (only background), when LR is negative, it implies that the events are more compatible with the alternative hypothesis H_1 (signal + background). Note that the variables θ' and θ'' of Equation 7.3 represent the nuisance parameters that maximize the likelihood for the two hypotheses. The next step is to define the confidence level, which is a measure of how reliable the experimental measurements established in the analysis are. The following equation gives the confidence value for the upper limit:

$$CL_s = \frac{CL_{s+b}}{CL_b} \quad (7.4)$$

Where CL_{s+b} quantifies how compatible the observed data is compared to the alternative hypothesis (signal + background) and CL_b quantifies the compatibility with the null hypothesis (background). The value of CL_s must be less than or equal to 0.05 to establish a confidence value of 95 % in the calculation of the limits.

All the previous procedure is executed automatically by using the Higgs limit tool [222]. This tool receives as input parameters some data cards that contain the yields and the nuisance parameters for each bin. Figure 7.5 shows the 95 % C.L upper limits obtained with the Higgs limit tool for six distinct values of $\Delta m = m_{\tilde{\chi}_1^\pm} - m_{\tilde{\chi}_1^0}$ (1, 10, 20, 30, 40, 50) for each leptonic channel. Figure 7.6 shows the results for the combination of all the channels.

In these figures, it is assumed that the decay of the smuon, selectron and stau can occur with the same probability (democratic scenario). Figure 1 shows that for values of $\Delta m = 1$ GeV the channel with the highest sensitivity is the invisible one, while for values of Δm above 10 GeV the channel that has the highest sensitivity is the muon channel.

Figure 7.7 shows the combined limits in a 2D plot for Δm vs $m(\tilde{\chi}_1^\pm)$ plane. The exclusion limit obtained for $\Delta m = 30$ GeV is 300 GeV, and for $\Delta m = 10$ GeV it is 212 GeV, which exceed the limits obtained in the previous searches at 8 TeV (<100 GeV).

7.5 Conclusions

- We performed a data analysis with final states of one muon, MET and two VBF jets on a data sample of 35.86 fb^{-1} collected by the CMS experiment in the year 2016. No data excess was observed over the expected SM background yields.
- The non-data excess result was interpreted under the conserving R-parity MSSM assuming electroweakino production through slepton decays with mass gaps between the chargino/neutralino and the LSP < 50 GeV.
- We obtained exclusion limits that extend the sensitivity in the compressed spectrum scenario:
 - The exclusion limit (95% CL) obtained for $\Delta m = 30$ GeV is 300 GeV.
 - For $\Delta m = 10$ GeV is 212 GeV.

These results surpass the limits reached in the previous searches at 8 TeV (<100 GeV).

- For values of Δm above 10 GeV, the channel that has the highest sensitivity for the s-lepton democratic scenarios is the muon channel in comparison with the other leptonic channels.
- For the high compressed spectrum scenario with $\Delta m = 1$ GeV, the channel with the highest sensitivity is the invisible channel. This is because the final state lepton p_T is very low.
- Even though we interpreted our results under a specific SUSY scenario, other models can also be tested with our results.
- We have proven that the signature with two VBF jets, MET, and one lepton is more sensitive than the two lepton final state because the leptons produced in the compressed scenario have soft p_T .

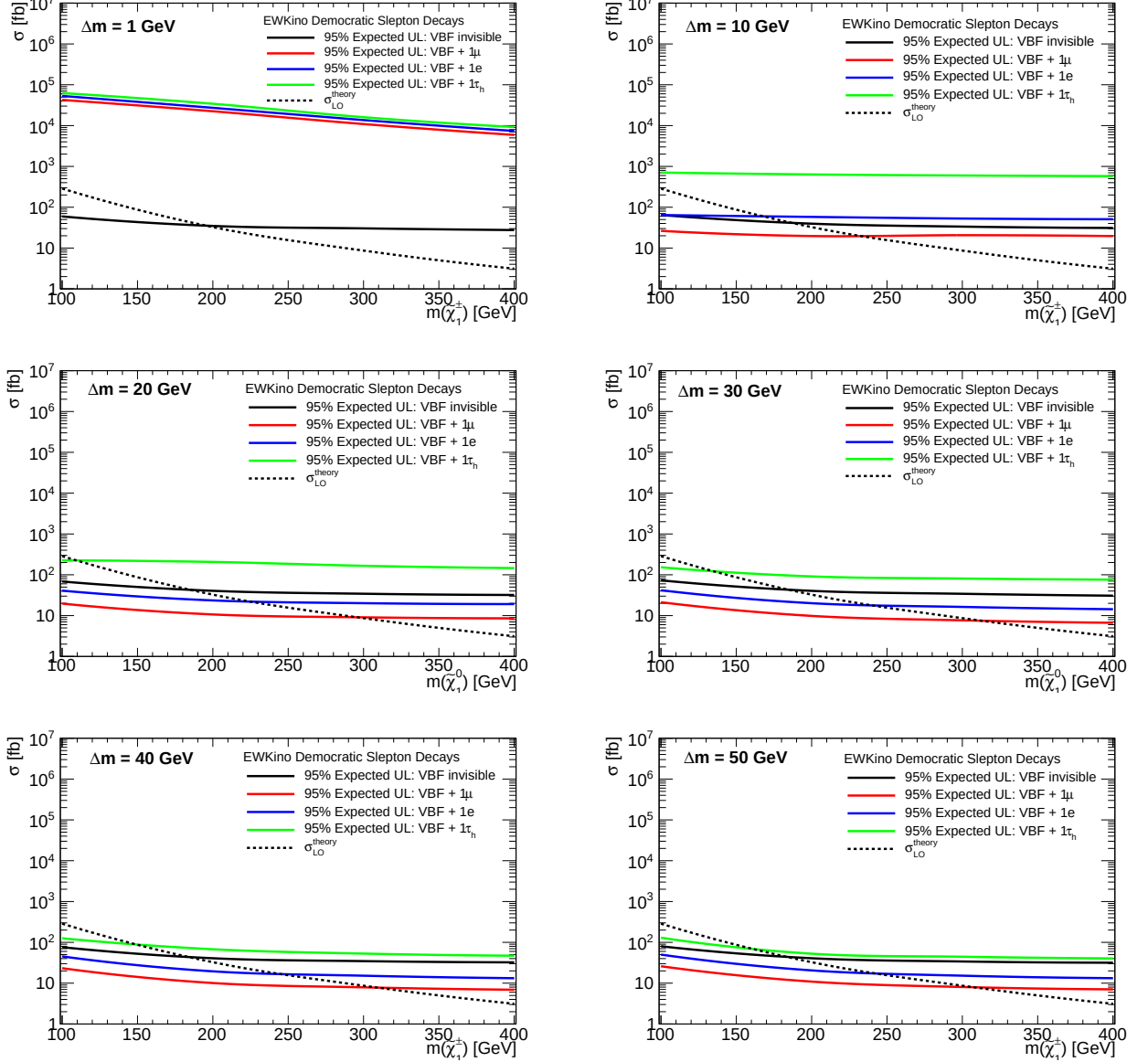


Figure 7.5: Upper limits for electrowikino signal cross section for VBF invisible search (black) and the single soft-lepton searches (red, blue, green), as a function of the chargino mass. The top row of plots correspond to $\Delta m = 1$ GeV (left) and $\Delta m = 10$ GeV (right) mass gap between the chargino and the neutralino. The middle row of plots correspond to $\Delta m = 20$ GeV (left) and $\Delta m = 30$ GeV (right). The bottom row of plots correspond to $\Delta m = 40$ GeV (left) and $\Delta m = 50$ GeV (right).

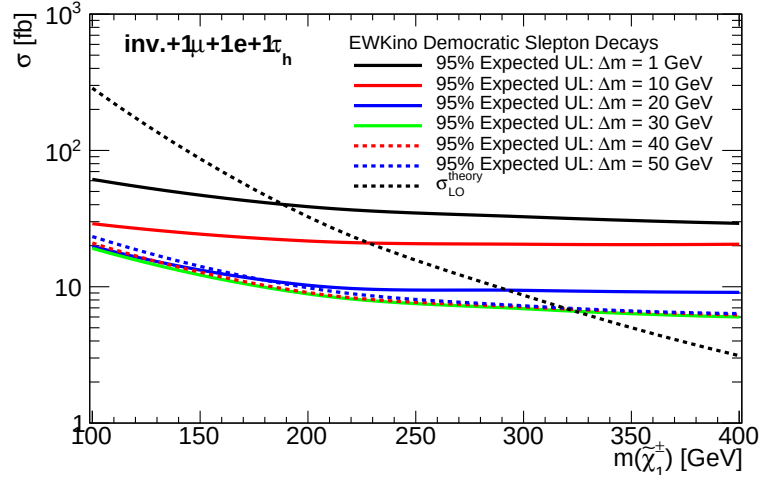


Figure 7.6: Combined limits.

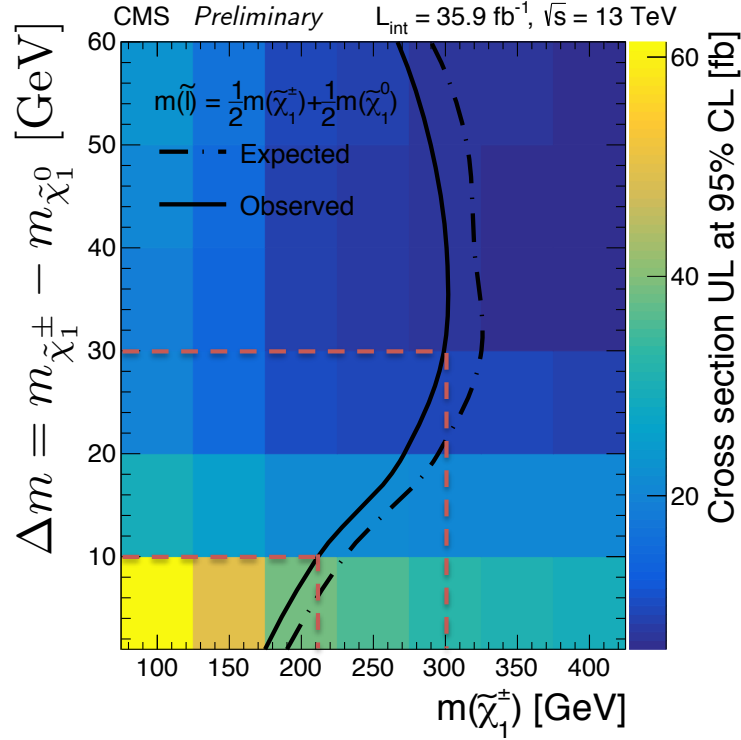


Figure 7.7: Combined limits in the Δm vs $m(\tilde{\chi}_1^\pm)$ plane. The exclusion limit obtained for $\Delta m = 30$ GeV is ~ 300 GeV, for $\Delta m = 10$ GeV it is ~ 212 GeV.

Table 7.3: MC predicted SM background contributions for the SUSY signal topology: VBF $jj + \ell + E_T^{\text{miss}}$, where $\ell=e, \text{tau}$.

VBF $jj + e + E_T^{\text{miss}}$ channel				
Sample	$N_{\text{BG}}^{\text{MC/Data}}(\text{SR})$	SF^{CR1}	SF^{CR2}	$N_{\text{BG}}^{\text{Data}}$
Diboson	1.1 ± 0.6	—	—	1.1 ± 0.6
QCD	$0^{+1.2}$	—	—	$0^{+1.2}$
Single Top	0.1 ± 0.1	—	—	0.1 ± 0.1
$W + \text{jets}$	5.3 ± 1.0	0.97 ± 0.10	1.18 ± 0.09	6.1 ± 1.4
DY + jets	0.1 ± 0.0	—	—	0.1 ± 0.0
$t\bar{t}$	11.4 ± 2.1	1.17 ± 0.09	0.82 ± 0.26	10.9 ± 4.1
BG Prediction	18.0 ± 2.4	—	—	18.3 ± 4.6
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 200, 195, 190 \text{ GeV}$	—	—	—	14.2 ± 0.4
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 200, 185, 170 \text{ GeV}$	—	—	—	37.3 ± 0.7
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 200, 175, 150 \text{ GeV}$	—	—	—	36.5 ± 0.7
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 300, 295, 290 \text{ GeV}$	—	—	—	3.4 ± 0.1
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 300, 285, 270 \text{ GeV}$	—	—	—	12.2 ± 0.2
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 300, 275, 250 \text{ GeV}$	—	—	—	13.1 ± 0.2
Data	—	—	—	29
VBF $jj + \tau_h + E_T^{\text{miss}}$ channel				
Sample	$N_{\text{BG}}^{\text{MC/Data}}(\text{SR})$	SF^{CR1}	SF^{CR2}	$N_{\text{BG}}^{\text{Data}}$
Diboson	0.5 ± 0.4	—	—	0.5 ± 0.4
QCD	619.9 ± 52.6	0.037 ± 0.008	—	23.1 ± 5.0
Single Top	0.5 ± 0.2	—	—	0.5 ± 0.2
$W + \text{jets}$	5.7 ± 1.1	1.04 ± 0.13	1.18 ± 0.09	7.0 ± 1.7
DY + jets	0.1 ± 0.1	—	—	0.1 ± 0.1
$t\bar{t}$	3.7 ± 1.2	1.06 ± 0.19	1.3 ± 0.5	5.1 ± 2.7
BG Prediction	630.4 ± 52.7	—	—	36.3 ± 5.9
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 200, 195, 190 \text{ GeV}$	—	—	—	0.38 ± 0.07
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 200, 185, 170 \text{ GeV}$	—	—	—	2.5 ± 0.2
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 200, 175, 150 \text{ GeV}$	—	—	—	4.3 ± 0.2
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 300, 295, 290 \text{ GeV}$	—	—	—	0.085 ± 0.02
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 300, 285, 270 \text{ GeV}$	—	—	—	0.74 ± 0.05
$m(\tilde{\chi}_1^\pm), m(\tilde{\ell}), m(\tilde{\chi}_1^0) = 300, 275, 250 \text{ GeV}$	—	—	—	1.4 ± 0.07
Data	—	—	—	38

Table 7.4: Signal Selection criteria for the invisible channel

Basic selection	Muon veto	$p_T(\mu) > 8 \text{ GeV}$ $ \eta(\mu) < 2.5$ Tight ID Isolation: $I < 0.25$
	Electron veto	$p_T(e) > 10 \text{ GeV}$ $ \eta(e) < 2.5$ Medium ID
	Tau veto	$p_T(\tau_h) > 20 \text{ GeV}$ $ \eta(\tau_h) < 2.5$ 1 prong $\Delta R(\tau_h, \mu \text{ or } e) > 0.3$
	Jets definition	$p_T(j) > 60 \text{ GeV}$ $ \eta(j) < 5$ Loose ID $N(j) \geq 2$
	b-tagged jets veto $N(b) = 0$	$p_T(b) > 30 \text{ GeV}$ $ \eta(b) < 2.4$ CSV Medium WP
p_T^{miss} criterion		$p_T^{\text{miss}} > 250 \text{ GeV}$
VBF		$\eta(j_1)\eta(j_2) < 0$ $ \Delta\eta(j_1, j_2) > 3.8$ $M(j_1, j_2) > 1 \text{ TeV}$
QCD rejection		$ \Delta\phi(p_T^{\text{miss}}, j_2) > 0.5$

Table 7.5: Summary of systematic uncertainties. Values are given in percentage, “s” indicates template variations (“shape” uncertainties).

Source	QCD	W	DY	$t\bar{t}$	VV	Signal
Lumi	2.5	2.5	2.5	2.5	2.5	2.5
μ ID	–	2	2	2	2	3
e ID	–	< 1	< 1	< 1	< 1	1
τ_h ID	–	1	1	1	1	2
Trigger	3	3	3	3	3	3
b ID	1	1	1	7	1	1
JES	s	s	s	s	s	s
TES	s	s	s	s	s	s
MMS	< 1	< 1	< 1	< 1	< 1	< 1
EES	< 1	< 1	< 1	< 1	< 1	< 1
pdf	–	4.8	4.2	4.2	3.5	6.5
scale	–	1	1	3.5	–	2
bin-by-bin stat.	s	s	s	s	s	s
Closure+Norm.	–	11	8	26	–	–

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