

Strong supersymmetry

A search for squarks and gluinos in hadronic channels
using the ATLAS detector

Robin van der Leeuw

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Introduction

High energy physics is the branch of science devoted to the study of the smallest, fundamental particles of the universe. The Standard Model of particle physics, constructed in the second half of the 20th century, describes all known fundamental particles and three of the four interactions between them; only gravity is not included. It has been extremely successful in describing physics at the smallest length scales, by making precise, verifiable predictions. Such predictions include the existence of the messenger particles of the weak interaction, the $W^\pm$ and Z bosons, which were discovered in 1983 at CERN in Geneva, and of the heaviest quark, the top quark, discovered in 1995 at FERMILAB in the US.

The last missing piece of the Standard Model puzzle was discovered at CERN in 2012: the Higgs boson. Introduced in the theory in the 1960s, it is part of a mechanism breaking the so-called electroweak symmetry, resulting in massive Standard Model particles. With the Higgs particle discovered, and the Standard Model complete, at first glance particle physics seems to have come to its conclusion. However, looking closer, there are several issues with the model.

The astronomical discovery of *dark matter*, which seems to have no or very weak interactions with ordinary matter except via gravity, has led to the belief that the Standard Model does not describe *all* the particles in the universe. Only 18% of all matter in the universe is normal Standard Model matter, with dark matter being much more abundant. The accelerating expansion of the universe points to a form of energy which cannot be accounted for by the current physical theories, dubbed *dark energy*. Furthermore, since if in the early universe equal amounts of matter and anti-matter would have been generated, it would have all annihilated, there must have been a matter-antimatter asymmetry. The current Standard Model does not lead to the required magnitude for the asymmetry.

There is also a theoretical issue with the model. Radiative corrections to the Higgs boson mass raise it to very high scales. The existence of the electroweak scale ($\mathcal{O}(100)$ GeV) at which the Higgs boson acts is therefore in some sense unnatural. This scale problem is known as the *hierarchy problem*.

Many theories have been devised to solve these issues, one of which is *supersymmetry*, which states that a symmetry exists in nature between fermions, particles with half-integer spin, and bosons, having integer spin. In supersymmetry, each known Standard Model particle has a partner particle with the same properties except for its spin, which differs by half a unit. Since none of these particles have been discovered so far, supersymmetric particles need to be heavier than their partners. This can be achieved by breaking the symmetry. Supersymmetry is currently the most researched

theory providing a solution to the hierarchy problem while being able to provide a candidate for a dark matter particle.

To test the Standard Model, particle accelerators are used, which accelerate particles and collide them at high centre-of-mass energies. From the relationship between energy and mass, $E = mc^2$ (made famous by Einstein), it becomes clear that the energy of two particles which collide with high velocity can be transformed into mass of heavier particles. However, the heavier the mass of produced particles, the higher the initial energy needs to be. To search for supersymmetric particles, the Large Hadron Collider (LHC) at CERN is used. The LHC accelerates protons to 99.9999% of the speed of light, and collides them with a centre-of-mass energy of 8 TeV¹. The decays of the particles produced in the collisions are detected by particle detectors, such as the ATLAS detector.

Previous experiments have not found evidence for supersymmetric particles, and set limits on the mass of supersymmetric particles. However, with the high collision energies achieved at the LHC the searches for these particles have entered a new realm of possibilities. This thesis describes a search for supersymmetry with the ATLAS detector. We search for supersymmetry by studying discrepancies from the Standard Model in data recorded by ATLAS.

Thesis outline

This thesis is organised as follows: a theoretical foundation is provided in chapter 1, which covers the theory of the Standard Model of particle physics in short, and subsequently describes several of its shortcomings. The second half of the chapter discusses the theory and phenomenology of supersymmetry, motivating the search for supersymmetry in LHC collisions. Chapter 2 describes the design and performance of the LHC and the ATLAS detector. The operating performance of the semiconductor tracker, one of the tracking components of ATLAS for which I was responsible for monitoring the efficiency, is treated in greater detail.

The rate of supersymmetric particle production is given by the production cross section, which is computed numerically with inputs from analytical calculations and experimental results, such as parton distribution functions. Both these components have uncertainties, which need to be propagated to uncertainties on the cross sections for supersymmetric particle production. I have been responsible for developing a framework to calculate these cross sections and their uncertainties for SUSY signal processes. This framework is used by the SUSY working group in ATLAS, and is described in chapter 3.

Particles traversing the detector are not observed directly, but need to be reconstructed from electronic detector signals. This event reconstruction is discussed in chapter 4, where also the data simulation using Monte Carlo techniques is described. If strongly interacting supersymmetric particles are present in nature with masses around 1 TeV, they should be abundantly produced in LHC collisions and subsequently decay into highly energetic quarks, possible leptons and undetectable stable particles.

¹1 TeV is equivalent to the energy of a flying mosquito, which consists of the order of 10^{23} atoms.

Therefore, we use the reconstructed events to search for these types of supersymmetric particles in events with energetic jets and high missing transverse momentum without electrons or muons. This search, using 5.8 fb^{-1} of $\sqrt{s} = 8 \text{ TeV}$ data, is discussed in detail in chapter 5. I have been working on this analysis and its predecessor, conducted on $\sqrt{s} = 7 \text{ TeV}$ data, since 2011. Within the group of ATLAS physicists collaborating on this analysis, I have been in charge of defining the analysis for certain supersymmetric scenarios with small mass differences between the supersymmetric particles, including the signal grid definition, analysis optimisation and setting the exclusion limits. Furthermore, I have worked on other areas of the analysis, such as the background estimation and the analysis framework development and maintenance.

Finally, in chapter 6 the results and implications of this analysis are discussed, and compared to other SUSY analyses done by ATLAS.

Theory of the Standard Model and Supersymmetry

Particle physics is the field of science which describes the fundamental building blocks of matter: subatomic particles, and the fundamental forces between them. Although the idea that matter consists of elementary particles dates back to the ancient Greeks, the modern view developed in the late nineteenth century and the first half of the twentieth century, with the discovery of the electron, the atomic nucleus and its structure of protons and neutrons, and the development of quantum mechanics. It developed into the Standard Model (SM) of particle physics [1–4] during the second half of the twentieth century, describing accurately all known particles together with three of the four fundamental forces – only the gravitational force is not part of the Standard Model.

In this chapter the theoretical framework of this thesis will be given. In the first section the Standard Model will be summarised, beginning with a short description of the particles and forces. The Standard Model Lagrangian will be introduced in section 1.1.1, where quantum electrodynamics and the electroweak theory are explained. The strong force, described by quantum chromodynamics, and its consequences, is described in more detail in section 1.1.2.

The second section goes shortly into the shortcomings of the Standard Model, while the third section describes a theory which addresses some of these shortcomings: supersymmetry.

1.1 The Standard Model

The Standard Model is a relativistic quantum field theory, describing elementary particles and the fundamental forces interacting between them. Here *elementary* and *fundamental* are keywords: elementary particles are not built of smaller constituent particles, while the fundamental forces cannot be reduced to more basic forces. From elementary particles larger, composite particles can be built, such as protons and neutrons, leading to even larger structures such as atoms and molecules. The Standard

Model contains a combined description of three of the four fundamental forces: the electromagnetic force, the weak and the strong nuclear force. The fourth fundamental force, gravity, has not been integrated in this particle theory, as the unification of general relativity with a particle description of gravity has been proven to lead to singularities. The electromagnetic force is responsible for most of the macroscopic phenomena encountered in daily life: apart from the obvious influences such as electromagnetic radiation (e.g. light) and magnetism, it is the driving force behind most chemistry processes. The weak force is responsible for the radioactive β decay, while the strong force binds protons and neutrons together to form atomic nuclei. Section 1.1.1 will go into the details of the quantum field theories behind these forces, yet first a short description of the particles and forces in the SM is given.

Unlike the previous notion of fundamental building blocks of nature, atoms, the elementary particles in the Standard Model are assumed to be point-like, i.e. without spatial dimension. In 1911 Ernest Rutherford discovered that an atom consists of a nucleus with electrons surrounding it [5]. By the early 1930s, it was shown that the nucleus consists of protons and neutrons by Chadwick [6], Ivanenko and others [7]. Yet also these were shown not to be elementary particles: in 1964, motivated by the discovery of over a hundred strongly interacting composite particles, Murray Gell-Mann [8] and George Zweig [9] independently proposed a model where neutrons and protons consist of quarks – the proton is made of two *up-quarks* and one *down-quark*, while a neutron has one up-quark and two down-quarks. It is believed that these quarks do not contain an underlying structure, and thus are elementary. Similarly theory states that the electron, already discovered in the late nineteenth century, is an elementary particle.

While protons, neutrons and electrons are the building blocks of all matter on earth, the Standard Model consists of more elementary particles than just the up-quark, down-quark and electron. The particle content can be divided into two groups, *fermions* and *bosons*. Fermions have half-integer spin and obey Fermi-Dirac statistics, while bosons have integer spin and obey Bose-Einstein statistics. In the Standard Model, the former are the matter particles, while the latter are force carrying particles, mediating the interactions between particles. Figure 1.1 shows an overview of the various particles. Besides these, each elementary particle has a corresponding anti-particle, with opposite charge. From here onwards, statements regarding particles hold for anti-particles as well, unless stated otherwise.

The fermions are subdivided in six quarks and six *leptons*, which are again categorized in three generations or families. The second and third generation particles have the same quantum numbers as the first, yet are heavier. The first generation consists of two quarks, up (u) and down (d), and two leptons, the electron (e) and electron-neutrino (ν_e). Quarks are electrically charged particles, subject to the electromagnetic, weak and strong force. There are three up-type quarks with an electric charge of $+2/3 e$ ¹: besides the first generation up quark, there are the *charm* (c) and *top* (t) quarks, of second and third generation, respectively. The down-type quarks

¹The electrical charge is given in units of the elementary charge, $e = 1.6 \times 10^{-19}$ C

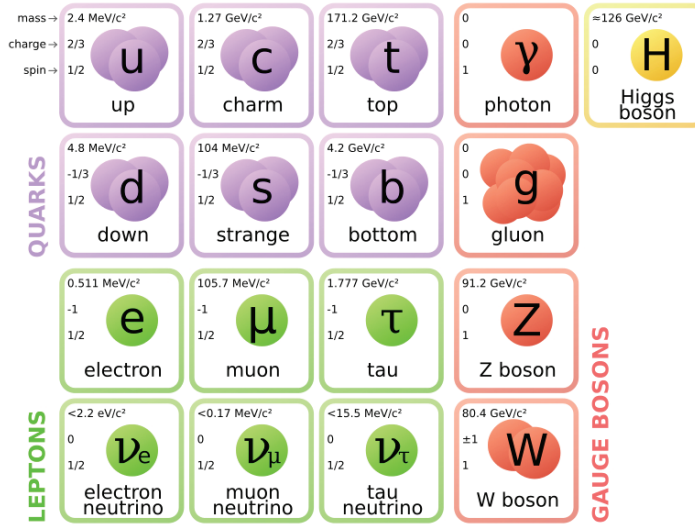


Figure 1.1: The particles of the Standard Model of particle physics. [10]

are, besides the down-quark, the *strange* (*s*) and *bottom* (*b*) quarks, which have an electrical charge of $-1/3 e$.

Besides the first generation electron (*e*) and electron-neutrino (ν_e), the second and third generation leptons are the *muon* (μ) and *tau* (τ), together with the muon- and tau-neutrino, respectively. The electron, muon and tau all have charge $-1 e$, while the neutrinos are neutral particles. This means charged leptons are subject to the electromagnetic and weak force, while the neutrinos interact only via the weak force – none of the leptons interact via the strong force.

The interactions between these matter particles are mediated through the gauge bosons. For the electromagnetic force, this is done by the photon (γ), which is a massless and neutral spin 1 particle. It couples to electrically charged particles, and thus does not interact with neutral particles. The weak force is mediated by three gauge bosons, $W^\pm$ and Z , which are massive. The weak force is a short-range force, acting on ranges of 10^{-16} – 10^{-17} m. It is the only interaction allowing for the decay of elementary particles, by being capable of changing the flavour of quarks and leptons. The strong force is mediated by eight massless gluons (*g*). Just like the electromagnetic force, the strong force has a conserved quantum number corresponding to it, called *colour charge*. Quarks are colour charged, while all leptons are colourless. Gluons are colour charged as well: they carry one of eight possible superpositions of a colour state and an anticolour state. As gluons only couple to colour charge, they do not interact with other bosons or leptons. Coloured particles have never been observed independently – they always come in combinations of three quarks (baryons) or a quark-antiquark pair (mesons), constructing colour singlet states with an integer electrical charge.

1.1.1 The Standard Model Lagrangian

The Standard Model is a quantum field theory, where each particle is described by a field in space-time. In quantum field theories, just like Newtonian physics, an equation of motion of Standard Model particles is obtained from a Lagrangian. As a viable theory of physics, it should most importantly meet the requirements that all observed particles (discussed above) should be included; it should be invariant under translations, rotations and boosts, i.e. under transformations of the Poincaré group; and it should be a renormalisable theory. The Standard Model is based on the $SU(3)_C \times SU(2)_L \times U(1)_Y$ symmetry group, where the strong interactions are described by the $SU(3)_C$ group, and the unified electroweak interactions are described by $SU(2)_L \times U(1)_Y$ – here the subscripts C and Y denote the colour charge and the hypercharge, quantum numbers conserved by their respective symmetry group. The subscript L indicates coupling to only left-handed fermions. It is assumed that the Standard Model is invariant under local transformations in this symmetry group, which has large implications.

The most straightforward Lagrangian density² possible for massless fermion fields ψ without interactions is

$$\mathcal{L} = \bar{\psi} i \gamma^\mu \partial_\mu \psi, \quad (1.1)$$

which contains just a kinetic term. Here γ^μ are Dirac matrices satisfying the anti-commutation rule $\{\gamma^\mu, \gamma^\nu\} = 2\eta^{\mu\nu}$, with $\eta^{\mu\nu}$ the Minkowski metric, ∂_μ is a partial derivative, and the repeating index μ implies a sum over all values of the index, running from 0 to 3. The fermion field ψ is a spinor, an element of $SU(2)$ which can be represented by a 4-component vector. To get to a more physical model, interactions between particles need to be described, be it either electromagnetic, weak or strong interactions. These are obtained by using gauge symmetries: changes to the fields under which the Lagrangian is invariant. For instance, the Lagrangian in equation 1.1 is invariant under a global transformation of the form $\psi(x) \rightarrow e^{i\alpha} \psi(x)$.

In the following, a short summary of quantum electrodynamics, the weak force, Higgs mechanism and quantum chromodynamics will be given, which will build up the Standard Model Lagrangian.

Quantum Electrodynamics

The electromagnetic interactions are described by the relativistic quantum field theory called quantum electrodynamics (QED), which is based on the symmetry group $U(1)$. The interaction between electrically charged particles is mediated by the massless neutral photon. The Lagrangian given in equation 1.1 can be updated to include a coupling between the fermions and the photon field, given by A^μ , by using the foundations of a local $U(1)$ symmetry: the Lagrangian should be invariant under a *local* gauge transformation of the form $\psi(x) \rightarrow e^{i\alpha(x)} \psi(x)$. Although equation 1.1 is not invariant under this transformation, it can be achieved by replacing the derivative

²Hereafter the Lagrangian density will be just called ‘Lagrangian’

by a covariant derivative

$$D_\mu = \partial_\mu - ieA_\mu, \quad (1.2)$$

where the vector potential A_μ transforms as

$$A_\mu \rightarrow A_\mu + \frac{1}{e}\partial_\mu\alpha(x). \quad (1.3)$$

Inserting the covariant derivative into equation 1.1 and adding a kinetic term for A_μ , which is gauge invariant, gives:

$$\mathcal{L}_{QED} = \bar{\psi}(i\gamma^\mu D_\mu - m)\psi - \frac{1}{4}F^{\mu\nu}F_{\mu\nu}. \quad (1.4)$$

$$F^{\mu\nu} = \partial^\mu A^\nu - \partial^\nu A^\mu. \quad (1.5)$$

This Lagrangian describes the fermion-photon coupling, with strength $\alpha = e^2/(4\pi\epsilon_0\hbar c) \approx 1/137$, if we take A_μ as the photon field.

Although seemingly very theoretical after such a short introduction, QED has many experimental predictions, and thus has been thoroughly tested experimentally. The most accurate measurement to date probing QED has been of the anomalous magnetic dipole moment of the electron, $(g_e - 2)/2$, which examines the interaction of electrons with fluctuations in the vacuum – the predicted value has been confirmed up to $|\delta(g_e - 2)/2| < 8 \times 10^{-12}$ [11], with $|\delta(g_e - 2)/2|$ the difference between the observed experimental value and the theoretical prediction.

The weak interaction and electroweak theory

After the discovery of the neutron by Chadwick in 1932 [6], and Pauli's postulate of the neutrino in 1930 [12] which provided a solution to the beta decay problem³, Fermi proposed a new theory of beta decay. Introducing a new force, this theory described how a neutron decays in a proton, electron and a neutrino [14]. The new force has a strength, given by Fermi's constant G_F , much lower than the electromagnetic coupling or the later discovered strong interaction, and is thus called the weak interaction. The 'charge' belonging to the interaction is the weak isospin: T , of which the third component, T_3 , is conserved under weak interactions, and is thus usually taken to be the weak isospin. Left-handed fermions, see below, have values of $T_3 = \pm 1/2$, while bosons have $T_3 \in \{\pm 1, 0\}$. Up-type quarks (up, charm, top) with $T_3 = +1/2$ can only transform via the weak interaction in down-type quarks (down, strange, bottom) which have $T_3 = -1/2$, and vice versa.

In 1968 Glashow, Weinberg and Salam [1–3] proposed a unified theory of electromagnetic and weak interactions, the electroweak theory: a quantum field theory with a

³In 1914 Chadwick observed that beta decay, i.e. the radiation of an electron by an atom, was emitted with a continuous energy spectrum, unlike alpha and gamma radiation. In the early 1920s beta decay was thought to occur as $N_i \rightarrow N_f + e^-$, with N_i, N_f the initial and final nucleus. This assumption leads to a discrete quantity for the kinetic energy of an electron – yet a continuous spectrum was observed, leading to a violation of energy conservation. For a detailed history, see [13].

$SU(2)_L \times U(1)_Y$ gauge symmetry. The subscript L denotes that only the left-handed fermions are subject to the weak interaction. The fermions are described by Dirac spinors, which can be decomposed in a left- and right-handed chirality eigenstate ψ_L and ψ_R ⁴:

$$\psi = P_L \psi + P_R \psi = \psi_L + \psi_R \quad (1.6)$$

$$P_L = \frac{1 - \gamma^5}{2}, \quad P_R = \frac{1 + \gamma^5}{2}. \quad (1.7)$$

Here P_L, P_R are the left- and right-handed projection operators, and the matrix operator γ^5 is given by $\gamma^5 = i\gamma^0\gamma^1\gamma^2\gamma^3$. Within the electroweak theory, the left-handed fermion fields appear as doublets: $L_i = \begin{pmatrix} \nu_i \\ l_i^- \end{pmatrix}$ for i th generation leptons and $Q_i = \begin{pmatrix} u_i \\ d_i \end{pmatrix}$ for i th generation quarks. The right-handed fermion fields (E_i for leptons, U_i, D_i for up and down type quarks) are singlets which do not couple to the weak interaction: their weak isospin is zero. On the other hand, left-handed doublets *do* couple to the weak interaction: under a $SU(2)$ transformation, weak isospin is conserved. In electroweak theory, weak isospin is combined with the electric charge Q to form the weak hypercharge $Y_W = 2(Q - T_3)$, which is the conserved quantity of the $U(1)_Y$ symmetry.

The Standard Model is built such that for each gauge symmetry, there are corresponding gauge fields. For the electroweak $SU(2) \times U(1)$ gauge symmetry these are the gauge fields W_μ^i corresponding to the $SU(2)$ symmetry, and a B^0 boson corresponding to $U(1)$. With these fields, a Lagrangian can be constructed analogous to the one of QED (equation 1.4):

$$\mathcal{L}_{EW} = i \sum_f \bar{\psi}_f \gamma^\mu D_\mu \psi_f - \frac{1}{4} W_{\mu\nu}^i W^{i\mu\nu} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu}, \quad (1.8)$$

where $W_{\mu\nu}^i$ and $B_{\mu\nu}$ are the field strength tensors for the fields W_μ^i and B_μ^0 , given by:

$$W_{\mu\nu}^i = \partial_\mu W_\nu^i - \partial_\nu W_\mu^i + g\epsilon_{ijk} W_\mu^j W_\nu^k \quad (1.9)$$

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu. \quad (1.10)$$

Here ϵ_{ijk} is the Levi-Civita tensor, which is $+1$ for even permutations of $(i, j, k) = (1, 2, 3)$, and -1 for odd permutations. Just like the QED Lagrangian, the electroweak Lagrangian in equation 1.8 is invariant under a local $SU(2) \times U(1)$ symmetry if the covariant derivative is chosen correctly:

$$D_\mu = \partial_\mu + \frac{1}{2}ig\tau^i W_\mu^i - \frac{1}{2}ig'Y B_\mu. \quad (1.11)$$

⁴The chirality of a particle determines if it transforms under a right or left-handed representation of the Poincaré group. In the relativistic limit, a particle's chirality equals its helicity: the spin direction of a particle with right-handed helicity is in the same direction as its motion; for left-handed particles it is in opposite direction.

Here g and g' are the coupling strengths of $SU(2)$ and $U(1)$, respectively, and the Pauli matrices τ^i and hypercharge Y are the generators of the $SU(2)$ and $U(1)$ symmetry, respectively.

Electroweak symmetry breaking and the Higgs mechanism

One major problem with the symmetry as stated above is that it leads to massless gauge bosons. The low strength of the weak interaction, and its limited range, can be explained by mediating bosons with mass – also, the lack of evidence for massless bosons pointed towards bosons with mass, years before the discovery of massive bosons. Yet adding a boson mass term to the Lagrangian leads to breaking gauge invariance. This was solved by introducing the Higgs mechanism⁵ of electroweak symmetry breaking (EWSB) [15–17].

EWSB is performed by adding a complex scalar $SU(2)$ doublet Φ made of two complex fields ϕ^0 and ϕ^+ , together containing four scalar degrees of freedom, $\phi_{1,2,3,4}$:

$$\Phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} \equiv \begin{pmatrix} \frac{1}{\sqrt{2}}(\phi_3 + i\phi_4) \\ \frac{1}{\sqrt{2}}(\phi_1 + i\phi_2) \end{pmatrix}. \quad (1.12)$$

The scalar doublet is added to the Lagrangian through the terms

$$\mathcal{L}_{Higgs} = D_\mu \Phi^\dagger D^\mu \Phi - \mu^2 (\Phi^\dagger \Phi) - \lambda (\Phi^\dagger \Phi)^2 \quad (1.13)$$

where the covariant derivative is again given by equation 1.11 to preserve the $SU(2) \times U(1)$ invariance, and the last two terms make up the Higgs potential $V_H(\Phi)$. The breaking of the electroweak symmetry is obtained by taking a negative mass parameter, $\mu^2 < 0$, while the self-coupling is non-zero and positive, $\lambda > 0$. The potential is thus a so-called ‘Mexican hat’, with a degenerate minimum of the potential at non-zero value, $|\Phi|^2 = -\mu^2/(2\lambda) \equiv v^2/2$. Here $v = \sqrt{-\mu^2/\lambda}$ is the vacuum expectation value (vev), which corresponds to the ground state of the field Φ . Gauge freedom allows for the choice of unitarity gauge for the field Φ , where both imaginary components (ϕ_2, ϕ_4) and one real component (ϕ_3) is set to zero, and we are left with only one real component (ϕ_1):

$$\Phi = \begin{pmatrix} 0 \\ \frac{v}{\sqrt{2}} \end{pmatrix} \quad (1.14)$$

Variations around the minimum of the potential are obtained by expanding Φ around the ground state: $\phi_1 = v + h$ with h a real scalar field. Thus the Lagrangian is written as

$$\begin{aligned} \mathcal{L}_{Higgs} &= \partial_\mu h \partial^\mu h - \lambda v^2 h^2 - \lambda v h^3 - \frac{\lambda}{4} h^4 \\ &+ \frac{1}{8} [g^2 (W_1^2 + W_2^2) + (-g W_\mu^3 + g' Y B_\mu)^2] (v + h)^2, \end{aligned} \quad (1.15)$$

⁵Although the name is under increasing debate, this thesis will use the name ‘Higgs mechanism’ in favour of ‘Anderson-Englert-Brout-Higgs-Guralnik-Hagen-Kibble mechanism’, which although historically and politically more correct, is just plainly too long.

where on the first line a mass term arises for the Higgs field of the form $m_h = \sqrt{2\lambda v^2}$, together with a three and four point self coupling. The second line introduces new couplings between the Higgs and vector fields, and mass terms for the vector bosons via a non-diagonal mass matrix. The matrix can be diagonalised by mixing the vector fields:

$$W_\mu^\pm = \frac{1}{\sqrt{2}}(W_\mu^1 \mp W_\mu^2) \quad (1.16)$$

$$\begin{pmatrix} A_\mu \\ Z_\mu \end{pmatrix} = \begin{pmatrix} \cos \theta_W & \sin \theta_W \\ -\sin \theta_W & \cos \theta_W \end{pmatrix} \begin{pmatrix} B_\mu \\ W_\mu^3 \end{pmatrix}, \quad (1.17)$$

where the weak mixing angle θ_W is introduced for convenience, as

$$\cos \theta_W = \frac{g}{\sqrt{g^2 + g'^2}}, \quad \sin \theta_W = \frac{g'}{\sqrt{g^2 + g'^2}}. \quad (1.18)$$

The last line of the Lagrangian in equation 1.15 can be rewritten as

$$\mathcal{L} = \frac{1}{2}g^2 W^+ W^- (v + h)^2 + \frac{1}{8}(g^2 + g'^2)Z^2(v + h)^2 \quad (1.19)$$

which does not contain any quadratic term in the field A_μ – the photon has not gained any mass. The masses of the W and Z boson and photon are thus:

$$m_W = gv/2 \quad (1.20)$$

$$m_Z = \frac{m_W}{\cos \theta_W} \quad (1.21)$$

$$m_A = 0. \quad (1.22)$$

By introducing the Higgs field, three of the four massless vector bosons have gained mass, while the photon field is still massless, just as required. The masses of the W and Z boson depend on v , which can be determined through its relation to the Fermi constant, $v = (\sqrt{2}G_F)^{-1/2}$, which is measured from the muon lifetime. Given the measurement of $v = 246.22$ GeV [18] the predicted masses⁶ agree well with the measured values of $m_W = 80.385 \pm 0.015$ GeV and $m_Z = 91.1876 \pm 0.0021$ GeV [19]. At the same time, a new scalar boson is added to the Standard Model. The discovery of this Higgs boson [20, 21] with a mass of $125.5 \pm 0.2(\text{stat})_{-0.6}^{+0.5}(\text{sys})$ GeV [22] has been the biggest achievement of the LHC to date.

The addition of the Higgs field also allows fermions to have mass. Fermion mass terms of the form $m_f(\bar{\psi}\psi) = m_f(\bar{\psi}_L\psi_R + \bar{\psi}_R\psi_L)$ were previously forbidden in the Standard Model as they violated $SU(2)$ symmetry. This is solved by introducing gauge invariant Yukawa couplings between the left-handed doublet, right-handed singlet and

⁶Throughout this thesis natural units are used, with $c = \hbar = 1$. Masses are thus given in eV instead of eV/c^2 .

the Higgs field: after EWSB, for leptons we have

$$Y_{e,i}\bar{L}E\phi = Y_{e,i}e_L e_R \frac{v+h}{\sqrt{2}} = m_{e,i}e_L e_R + \frac{Y_{e,i}}{\sqrt{2}}e_L e_R h, \quad (1.23)$$

with i the generation, and likewise for quarks. The first term is the mass term coupling left- and right-handed fermions, with $m_{e,i} = Y_{e,i}v/\sqrt{2}$, while the second is the Higgs coupling to them.

One omission in the electroweak theory of the Standard model is neutrino mass. Although no right-handed counterpart exist for neutrinos in the Standard Model, and therefore no mass term either, the discovery of neutrino oscillations indicates neutrinos do have finite, small mass [23]. Solutions include the possibility that the neutrino is a Majorana fermion, for which the neutrino and anti-neutrino are the same particle; and sterile neutrinos, for which the right-handed neutrino would be an $SU(2)$ -singlet which only couples to gravity.

1.1.2 Quantum Chromodynamics

The third elementary force in the Standard Model is the strong interaction, which is described by quantum chromodynamics (QCD). The strong interaction acts between coloured particles (quarks), and is mediated by gluons. As the LHC accelerates and collides protons, consisting of coloured quarks, QCD plays an important role in LHC physics. This section will give a more detailed description of the strong interaction and related topics, and will be a guide for several of the following chapters.

Like the other interactions, the strong interaction arises from local gauge invariance under a transformation of a symmetry group, in this case the non-abelian group $SU(3)_C$. This group has eight generators, corresponding to eight gluons, while the charge of the symmetry (colour) has three possible values: red, green and blue. The Lagrangian of QCD, for quark fields ψ_q and gluon fields G_μ^a , is given by

$$\mathcal{L}_{QCD} = \sum_q \bar{\psi}_{q,i}(i\gamma^\mu(D_\mu)_{ij} - m_q\delta_{ij})\psi_{q,j} - \frac{1}{4}G_{\mu\nu}^a G^{a\mu\nu}, \quad (1.24)$$

where the sum is over the quark flavours q and the implicit sum over the indices i, j is taken over the three colour charges in the fundamental representation of $SU(3)$. The form shows similarities with the QED Lagrangian, with a covariant derivative D_μ and a term quadratic in the field strength $G_{\mu\nu}$, which are given by

$$D_\mu = \partial_\mu + ig_s t^a G_\mu^a \quad (1.25)$$

$$G_{\mu\nu}^a = \partial_\mu G_\nu^a - \partial_\nu G_\mu^a - g_s f^{abc} G_\mu^b G_\nu^c. \quad (1.26)$$

The gluon field G_μ^a is given in the adjoint representation of $SU(3)$, with indices a, b, c running from 1 to 8 and coupling g_s . The generators of $SU(3)$ under this representation are given by the matrices t^a , which satisfy the commutation relation $[t^a, t^b] = if^{abc}t^c$. These gluons transform as octets under $SU(3)_C$, yet as singlets

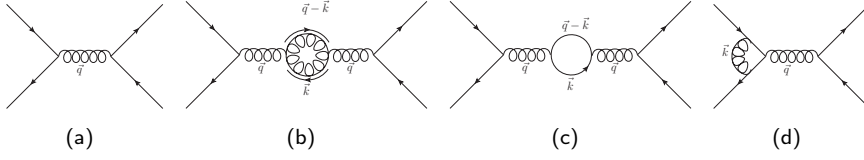


Figure 1.2: A leading order tree level diagram of quark-antiquark scattering (a), where a gluon with momentum $\vec{q}$ is exchanged. One-loop corrections to the gluon propagator can appear via a virtual gluon (b) or quark (c) loop with momentum $\vec{k}$. The exchange of a virtual gluon is also possible between the incoming quark and antiquark (d).

under $U(1)_Y$ and $SU(2)_L$, and thus only couple via the strong force. Unlike in QED, the mediating particles in QCD are colour charged themselves, making self interactions possible between the gluons: the last term in the Lagrangian allows for three- and four-point couplings of gluons. The coupling constant of QCD is given by $\alpha_s = g_s^2/(4\pi)$.

Two important properties of the strong interaction are *asymptotic freedom* [24, 25] and *colour confinement*, or just confinement. Asymptotic freedom states that the strong coupling strength α_s decreases with decreasing distance, or higher energies: $\alpha_s \rightarrow 0$ for $Q^2 \rightarrow \infty$, with Q^2 the energy scale of the interaction. For low energies α_s is too large for perturbative expansions to be applicable, and thus analytical computations are not possible for so-called soft QCD. Yet due to asymptotic freedom perturbative expansions are possible for higher energies, or smaller distances. Confinement states that coloured particles cannot exist freely, and will always be confined in colour-neutral bound states, called *hadrons*, of which the proton is an example. When quarks are pulled apart, confinement dictates that a quark-antiquark pair is created out of the vacuum to form a new colour neutral state. In the LHC, this leads to *hadronisation* of high-momentum quarks into jets. Both the proton structure and hadronisation will be discussed further in this section.

Renormalisation

Within quantum field theories, the *cross section* of a process, i.e. the probability of a specific process to happen, is proportional to the square of the scattering amplitude. Within perturbative theories, such as the perturbative regime of QCD, these can be calculated using Feynman diagrams [26]. This leads to a power series in a small variable, yielding an infinite number of terms. Thus a cross section cannot be calculated exactly, but its calculation has to be done up to a certain order in the perturbative constant – in the case of QCD, this perturbative constant is α_s . As an example, take the simplest possible QCD interaction: quark-antiquark interaction. One of the leading order Feynman diagrams is shown in figure 1.2 (a). Yet to calculate the cross section of the interaction to more precision, one needs to include diagrams with quark and gluon loops, of which three are shown in figures 1.2 (b), (c) and (d).

A full description of quantum field theory is outside the scope of this thesis, yet

some topics which are important for this thesis will be explained in short, such as renormalisation and factorisation. For a comprehensive introduction into quantum field theories, see [27, 28].

The loop diagrams shown in figure 1.2 (b), (c) and (d) are actually quite problematic: the expressions for the diagrams include integrals over the momentum of the gluon or quark in the loop, which are of the form $\int_0^\infty d^4k$ and thus diverge. And not only does this happen once: the cross section includes an infinite sum of these divergent integrals, while the result must be a finite number which we can measure. To solve this, one needs *renormalisation*, introduced in two steps.

First, a prescription is needed to describe and tame the infinities: *regularisation*. Several regularisation schemes exist, one of which being the momentum cutoff scheme. In this scheme, a cutoff energy Λ is defined, which acts as an upper limit on momentum in the integrals, such that we do not integrate over all possible momenta. This leads to an outcome which depends on Λ .

The second step in cancelling these divergences is the renormalisation itself. The integrals are split: a renormalisation scale μ_R is introduced up to which the integral yields a finite result, while the divergent part, or Λ dependency, is removed by introducing counter terms. Although the introduction of counter terms seems arbitrary, it can be shown [27] that these counter terms can be incorporated in the fields and parameters of the theory, such as the coupling constants using the renormalisation group equation (RGE). In this way, the divergences are absorbed in the couplings, which now depend on μ_R , giving a Λ independent result. Note that introducing μ_R cannot be done without care: physical observables should not depend on a non-physical scale, only unphysical parameters, such as the coupling constant, can depend on μ_R . In 1972, Gerard 't Hooft and Martinus Veltman proved that all Yang-Mills theories (non-abelian gauge theories, such as the Standard Model) are renormalisable [4]: Standard Model observables are independent of μ_R when the infinite sum over all terms is performed.

However, practically speaking, in the calculation of observables such as cross sections, only a finite number of terms can be calculated, which each *do* depend on μ_R . The renormalisation scale is usually chosen to be equal to the energy scale of the interaction, as this choice eliminates large logarithms in the loop diagrams, and therefore optimises the convergence of the perturbative expansion [29]. The dependency on the choice of the renormalisation scale is therefore an uncertainty on these calculations, as will be shown further in chapter 3. For a more detailed discussion on regularisation and renormalisation, see for instance [27, 30].

The parton model and factorisation

As the LHC accelerates and collides protons, a correct description of protons and their constituents is imperative. Like all baryons, protons consist of three valence quarks, in this case two up quarks and one down quark. Due to quantum fluctuations, gluons and quark-antiquark pairs (also called *sea quarks*) can be briefly created before disappearing again, leading to more than three elementary particles per proton at any given time. All the constituents of the protons, i.e. the valence and sea quarks, together with the

gluons, are known as *partons*. This nomenclature derives from the first observation of elementary particles inside protons in deep inelastic scattering (DIS) experiments, which were only later known to be quarks and gluons. The proton's momentum is distributed among its constituent partons: every parton i carries a fraction x_i of the proton's momentum. This complicates the picture of hard scattering in the LHC: instead of just two particles with a specific momentum interacting during a collision, two unknown partons interact, both with an unknown momentum. Therefore it is customary to use a hadronic cross section: the probability of a specific interaction occurring between two hadrons. The relation between the hadronic cross section and partonic cross section is given by

$$\sigma_{hadronic} = \sum_{i,j} \int dx_1 dx_2 f_i(x_i) f_j(x_j) \sigma_{partonic}(x_i, x_j), \quad (1.27)$$

with a sum over gluons and all flavours of quarks, $x_{i,j}$ are the partonic momentum fractions and the integral runs over all allowed momentum fractions. The functions $f_k(x_k)$ are *parton distribution functions* (PDFs), giving the probability for finding a parton of type k with momentum fraction x_k inside the proton. In other words, the interaction between two hadrons is factored into a partonic hard interaction, which is a short distance effect, and the long distance behaviour of the partons inside the proton. This is known as *factorisation*. Although a full derivation of factorisation is outside the scope of this thesis, it can be explained in an intuitive manner when considering DIS, the scattering of an electron on a hadron at high momentum transfer, as nicely explained in ref. [29]. The given reasoning can be generalised to other inclusive cross sections. For a more detailed and quantitative description, see for instance [27, 29, 31].

Seen in the centre-of-mass system of the electron-hadron interaction, the hadron is Lorentz contracted due to its high energy (and thus speed), shortening the time it takes the electron to traverse the hadron. At the same time, with increasing energy the interaction time of the hadron's internal interactions are lengthened by time dilation, increasing the lifetime of virtual partons inside the hadron. When the time it takes the electron to traverse the hadron is smaller than the lifetime of virtual partons, an electron will interact with the hadron as a static state, with a definite number of partons. Electron-hadron interactions with a high momentum transfer will be mediated by a short-lived photon, which only interacts with one parton assuming the parton density is not too high. Thus the electron-hadron scattering can be interpreted as a scattering on just one parton in the hadron, carrying a momentum fraction x of the hadron's momentum in the centre-of-mass frame. Assuming no particles travel in the opposite direction to the hadron, x should be taken between 0 and 1.

Now the complicated high energy scattering has simplified, where the interactions between the partons can no longer interfere with the interactions between the electron and the parton. The total cross section of the scattering between an electron and a proton with momentum transfer Q^2 can now be defined using the above defined parton distribution function. It can be shown that for a value $x > \xi \equiv 2p \cdot q / Q^2$, with

p and q the momentum of the incoming hadron and virtual photon, respectively, the partons in the hadron are indeed free [29]. Now if the cross section for an electron with a parton is given by the Born cross section $\sigma_B(Q^2, x)$, the total DIS cross section is given by

$$\sigma_{eH}(\xi, Q^2) = \sum_i \int_{\xi}^1 dx f_i(x) \sigma_B(Q^2, x), \quad (1.28)$$

which is precisely of the form of equation 1.27. Note that the Lorentz contraction of hadrons prohibits partons from different hadrons to overlap and interact, modifying the parton distributions. This Lorentz contraction is thus paramount to the universality of the PDFs.

Yet something is missing from this reasoning. QCD states that the interacting quarks can radiate off gluons. Even more, the probability of emitting a (collinear) gluon increases for decreasing momentum of the radiated gluon, resulting again in divergent integrals. Yet as we have seen in renormalisation, divergent integrals can be dealt with by separating the divergent part from the finite terms at a factorisation scale μ_F . The divergent terms can be absorbed into the PDFs using the factorisation theorem, leading to an infrared safe expression of the partonic cross section, which can still be calculated perturbatively. The expression for the hadronic cross section becomes:

$$\sigma_{hadronic} = \sum_{i,j} \int dx_1 dx_2 f_i(x_i, \mu_F) f_j(x_j, \mu_F) \sigma_{partonic}(x_i, x_j, \mu_F). \quad (1.29)$$

This cross section itself is independent of the introduced factorisation scale, as required.

Parton distribution functions

Parton distributions are *universal*, i.e. they do not depend on the specific process for which one wants to know the cross section. This makes them measurable via precision measurements of well-understood processes, and then transferable to any other process. These measurements include measurements of deep inelastic scattering cross sections and W , Z and jet production at the TEVATRON. These measurements show that the number of partons in a proton depends on the squared energy scale of the scattering, Q^2 . At low Q^2 the valence quarks are more dominant in the proton, while for high Q^2 an increasing number of quark-antiquark pairs are visible in the proton with a low momentum fraction x . Figure 1.3 shows PDFs obtained by the MSTW collaboration for $Q^2 = 10 \text{ GeV}^2$ (left) and $Q^2 = 10^4 \text{ GeV}^2$ (right) as a function of x . An important observation is that actually gluons carry up to half of the proton momentum, with (anti)quarks only carrying the remaining half. The fraction carried by gluons even increases with rising Q^2 .

Many collaborations determine PDFs from fits to experimental data, the most common being MSTW [32], CTEQ [33] and NNPDF [34]. These all base their determination on the same procedure: a parametrisation of the PDFs at low Q^2 is made, where the unknown perturbative terms are negligible, either from assumptions or using a neu-

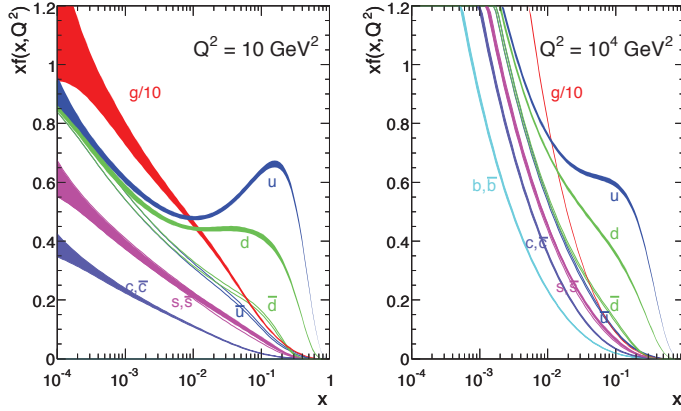


Figure 1.3: The distribution of $xf(x, Q^2)$, where x is the partonic momentum and $f(x, Q^2)$ are MSTW2008 PDFs, for an energy scale of $Q^2 = 10 \text{ GeV}^2$ on the left and $Q^2 = 10000 \text{ GeV}^2$ on the right. The PDFs are shown for the five lightest quarks and the gluon, where the band denotes the uncertainty [32].

ral network. Using the Dokshitzer-Gribov-Lipatov-Altarelli-Parisi (DGLAP) [35–37] evolution equations, these can be scaled to higher values of Q^2 , where a fit is performed on experimental data of e.g. cross section measurements. This data comes from various sources. First of all, DIS cross section measurements are used, both from fixed target lepton-nucleon experiments (electrons, muons and neutrinos scatter off hydrogen, deuterium and nuclear targets), which constrains quark and gluon PDFs at high x , and from electron-proton collisions at HERA which constrains them at low x . Yet DIS measurements mostly determine information on the valence quarks. To probe sea (anti)quarks, other datasets are used as well, such as fixed target Drell-Yan⁷ data to constrain high x sea quarks, single jet inclusive production cross section at the TEVATRON contributing to the high x gluon PDFs, and the W [38–40] and Z [41] boson production cross sections at the TEVATRON, which are sensitive to up and down quark distributions, and their anti-quark counterparts. The PDFs coming from these fits will depend on a large number of parameters, which, apart from the assumptions about the PDFs, originate from the choice of dataset, the specifics of the perturbative QCD calculation, the correlation between α_s and the PDFs, the treatment of the heavy quarks and lastly the uncertainty treatment. As each group differs on each of these points, the resulting PDFs will differ as well.

In this thesis PDF sets from the CTEQ and MSTW collaboration are used (CTEQ6.6 and MSTW2008). CTEQ PDFs are obtained from the above experimental data⁸ by

⁷In a Drell-Yan process between two hadrons, a quark and anti-quark annihilate to produce a Z boson or virtual photon (γ^*) decaying into two charged leptons.

⁸The most recent HERA 1 combined dataset is only included in the most recent CT10 PDF set

a global analysis [33]. These PDF sets are available for next-to-leading order (NLO) and next-to-next-to-leading order (NNLO) in perturbative QCD. The MSTW collaboration uses the same set of data for the determination of their PDFs, but performs a global analysis within the framework of leading-twist fixed-order collinear factorisation in the $\overline{MS}$ scheme [32]. It has PDFs for leading order (LO), NLO and NNLO QCD calculations.

Each of the parameters used in the fit to experimental data will have a corresponding uncertainty. The minimisation techniques used by the collaborations, usually based on the Hessian method [42], yield a best fit value for each of the parameters, with a corresponding χ^2 distribution. For every parameter of the CTEQ and MSTW PDFs the deviation is determined given a certain tolerance in the change in χ^2 . This tolerance in the different parameters is chosen to be $\Delta\chi^2 < 100$ (50) for CTEQ (MSTW), which results in a set of eigenvectors which describe the complete parameter set including the 90% (68%) confidence ranges. These sets of eigenvectors are used in chapter 3 to determine the uncertainties in SUSY cross section calculations due to PDFs and factorisation/renormalisation scales.

Hard scattering at the LHC

The total interaction between two protons colliding in ATLAS is a chaotic scene. Where the experiments at LEP had two leptons colliding, the protons in the LHC make calculations more difficult. Figure 1.4 shows an illustration of such a collision: the three valence quarks making up the incoming protons are seen left and right, with all stages of an interaction between two partons and the subsequent events shown.

The *hard scattering* process is the interaction between the incoming partons of the proton occurring with a large momentum transfer Q^2 , which can e.g. lead to energetic quarks in the detector, or the production of heavy particles such as top quarks. The resulting partons lose energy through *parton showering*, where soft and collinear quarks and gluons are radiated off until confinement takes over. At this point hadronisation takes place, recombining the partons into colour neutral hadrons. These hadrons will subsequently *decay* themselves into other hadrons, neutrinos, and occasionally leptons, resulting in a shower of particles called a *jet*. Apart from the hard scattering, a hard gluon can radiate off the incoming partons before or after the hard scattering takes place: *initial* or *final state radiation* (ISR or FSR). Both result in a jet which can end up in the detector. ISR jets are an important issue in the inclusive search for supersymmetry, as it provides additional jets to the hard scattering, as will become apparent in chapter 5.

The proton remnants which did not participate in the hard scattering are now coloured states. Due to confinement these states will interact via mostly soft scatterings with each other and hard scattering remnants to form colour neutral states. This hadronisation leads typically to soft jets along in the direction of the beam, yet can also enter the detector. This is called the *underlying event*.

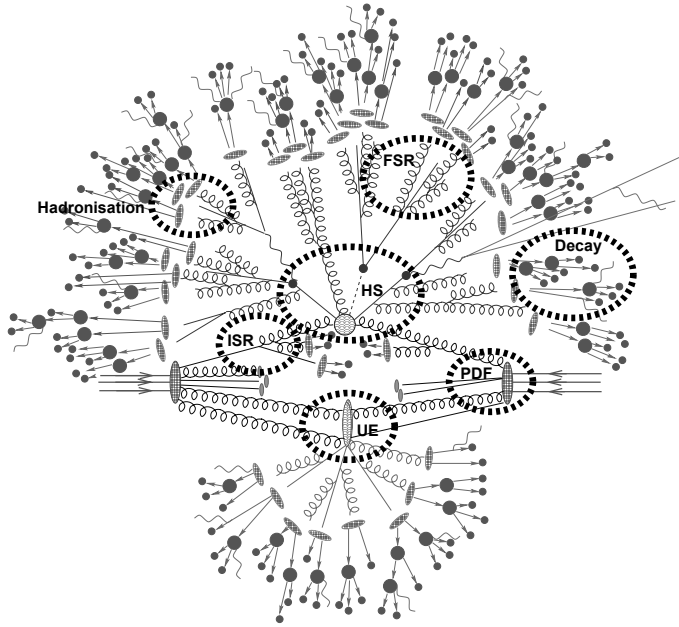


Figure 1.4: *Illustration of a general hard scattering event. All possible stages of the event are shown: the hard scattering (HS) from the incoming protons with their three valence quarks and corresponding PDFs; the initial state radiation (ISR) and final state radiation (FSR); parton showers with the subsequent hadronisation and decay into jets. The underlying event (UE) is shown in the lower half. Taken from [43] with modifications by [44].*

1.2 Shortcomings of the Standard Model: motivations for SUSY

The Standard Model is in a strange situation. On the one hand, it has been verified to high precision in past decades, with experiments verifying its predictions to exceptional precision. Yet on the other hand, there are various fundamental issues with the theory as will be shown in the next section, which lead to the search for evidence of physics *beyond the Standard Model*. Many theories have been devised to solve these issues, such as theories using multiple extra dimensions, technicolour or the little Higgs model. The theory which might be the most favoured by theorists, and is for sure the theory with the most attention from experimental physicists, is *supersymmetry*, or SUSY in short. As a motivation for a search for supersymmetry, several of the issues of the Standard Model are described below, together with solutions SUSY provides, after which SUSY theory is described in short.

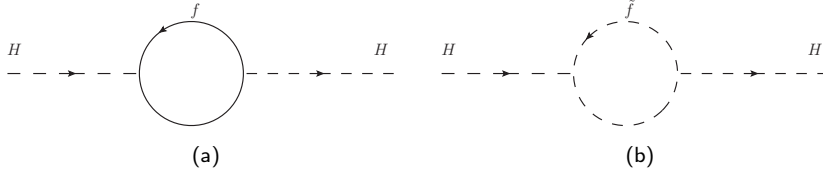


Figure 1.5: One-loop contributions to the Higgs self-energy with (a) a virtual Standard Model fermion f and (b) a supersymmetric scalar partner of this fermion, $\tilde{f}$. Dashed lines denote scalar particles, while solid lines denote fermions.

Hierarchy problem

One of the problems in Standard Model theory is known as the hierarchy problem. It is a conflict between energy scales in particle physics: on the one hand there is a requirement on the Higgs mass to be low, and on the other hand there are large loop contributions which drive the Higgs mass towards the largest energy scale available. Before its discovery, the Higgs mass was expected to be of order 100 GeV from precision measurements on weak interactions, while in order to prevent unitarity violation in gauge boson scattering it is required to be at least below 800 GeV [45], which are both in agreement with the mass of the discovered particle of $m_h \approx 126$ GeV. Yet when calculating the Higgs mass, one runs into trouble. Like all particles, scalars receive quantum corrections to their mass from particles which couple to it via virtual loop diagrams, such as shown in figure 1.5 (a) for a fermion f . This fermion with mass m_f couples to the Higgs with a coupling λ_f . The correction to the Higgs mass from these terms is proportional to

$$\delta m_h^2 \sim \lambda_f^2 \int d^4k \left(\frac{1}{k^2 - m_f^2} + \frac{m_f^2}{(k^2 - m_f^2)^2} \right). \quad (1.30)$$

The largest correction to the Higgs mass occurs when $\lambda_f \approx 1$, which is the case for top quarks. The first term in equation 1.30 is clearly quadratically divergent, while the second term has a logarithmic divergence. Introducing a cutoff scale Λ as explained in 1.1.2, the above integral results in $\delta m_h^2 \sim \lambda_f^2 \Lambda^2 + \mathcal{O}(\ln(\Lambda))$. According to the Standard Model, no new physics arises between the electroweak scale and the Planck scale, where quantum gravity effects become strong and some new physics is expected. The cutoff scale is therefore naturally taken as the Planck scale, $\Lambda = M_{Pl} = 1.2 \times 10^{19}$ GeV. As the Standard Model has been shown to be renormalisable, such quadratic divergences can be solved by introducing counter terms to the Lagrangian. Yet these counter terms should then remove corrections to the bare mass of the Higgs up to 17 orders of magnitude. Most theorists agree that such a level of fine-tuning of the model, although theoretically allowed, does not make for an aesthetically pleasing theory. And it is even a larger effect: the enormous amount of fine-tuning does not only affect the Higgs mass, but also that of every other Standard Model particle, as

it is really the Higgs mass parameter μ which is affected. The hierarchy problem as described here is therefore not so much a problem with the mass scales, as more a problem with fine-tuning. Note that although all fermion and gauge boson mass terms have divergent loop corrections, due to it being a scalar, the Higgs mass is the only quadratically divergent correction occurring in the Standard Model – corrections to other mass terms are at most logarithmically divergent, leading to much smaller, manageable corrections.

The quadratically divergent terms in the Higgs mass can be cancelled if scalar particles exist with the same quantum numbers as the fermions: the contributions to the Higgs mass would be equal, except for a minus sign coming from the fermion loop. Supersymmetry provides just that: it states that for each particle there exists a partner particle with the same mass and quantum numbers, but only differing in spin by $1/2$ unit. Figure 1.5 (b) shows such a virtual scalar partner of the fermion f contributing to the self-energy of the Higgs boson. The diagram shown will cancel exactly the quadratic divergence encountered in the Standard Model, as all couplings are the same. Likewise for fermionic partners of Standard Model bosons, which cancel bosonic contributions to the Higgs mass.

An apparent problem is that supersymmetry is seen to be broken, leading to a mass difference between the fermions and partner scalars (discussed in the next section). Yet if the mass difference is not too large, there still exists an approximate cancellation, needing considerable less fine-tuning. The partners of gluons and quarks, called gluinos and squarks, respectively, should have a mass of maximally a few TeV to remain within a fine-tuning up to 1 percent, which is usually considered to be the maximum "reasonable" fine-tuning.

Dark matter

Already since the 1930s astrophysical observations have indicated that there is much more matter than we can observe in the universe. Measurements on the orbital velocity of stars in our galaxy [46], and of galaxies in a cluster [47] indicated that more mass was needed to explain the orbits and keep our galaxy intact than is visible. The unknown matter is aptly named *dark matter*, and can only interact via gravity and the weak force. It thus cannot be electrically charged.

Observations done in the last 40 years are more compelling. When measuring the velocities of stars throughout galaxies, for a galaxy filled with 'normal' matter in stars as gas clouds, one would expect the velocity to decrease for large distance from the centre. However the rotation curve actually stabilises [48]. This can be explained if a halo of dark matter is present in the galaxy. Furthermore, WMAP [49] and more recently PLANCK [50] performed measurements on the temperature fluctuations of the cosmic microwave background radiation. When fitted to the current Standard Model of cosmology (Λ CDM), which describes a flat universe dominated by dark energy and dark matter, PLANCK's results show that only 4.9% of all energy in the universe is in the form of ordinary Standard Model matter, while 26.8% of all energy in the universe is of the form of dark matter and 68.3% is made up of dark energy [51].

The Standard Model does not contain a dark matter candidate, nor does it describe

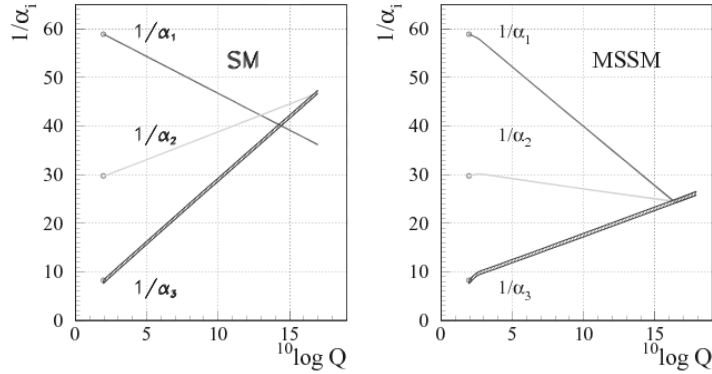


Figure 1.6: Evolution of the inverse of the coupling constants strength as a function of the energy Q for the Standard Model (left) and a Minimally Supersymmetric Standard Model theory with ~ 1 TeV particles (right). Here α_1 , α_2 and α_3 denote the electromagnetic, weak and strong couplings, respectively. Figure from [52].

a source for dark energy. SUSY on the other hand provides a natural candidate for dark matter: the lightest supersymmetric particle (LSP) is in many scenarios a stable, neutral and colourless particle, exactly what is needed for dark matter. Alas, the origin of dark energy is not explained by SUSY.

Unification of gauge couplings

A third issue to raise is not so much a problem of the Standard Model, as more a quest for a deeper theory. Seeing how two out of four fundamental forces have been unified into the electroweak force, the unification of all forces has been a holy grail in theoretical physics since many decades. A first step is to unify the electromagnetic, weak and strong force into one interaction via a Grand Unifying Theory (GUT), by embedding the $SU(3) \times SU(2) \times U(1)$ symmetry of the Standard Model in a larger symmetry group, such as $SU(5)$ or $E(6)$, which is broken at some higher energy scale. The Standard Model would then be a low energy effective theory.

Using the RGEs the evolution of the gauge couplings of the three fundamental forces can be described with increasing renormalisation scale. At an energy of $\Lambda_{GUT} \sim 10^{15}$ GeV, called the GUT scale, the strengths of the three forces become near identical, yet not exactly. This is shown in the left in figure 1.6. Additional heavy particles will influence this running of the gauge couplings, creating an opportunity for gauge coupling unification. When inserting TeV-scale SUSY particles to the theory, this is exactly what may happen. The right in figure 1.6 shows the renormalisation group evolution of the inverse gauge couplings assuming a SUSY theory with a SUSY breaking scale of ~ 1 TeV. The exact running of the couplings depends on the specific SUSY scenario, and on the masses of the particles – not all SUSY breaking mechanisms lead to unification, while the SUSY breaking scale must not be too high to achieve

unification.

Origin of electroweak symmetry breaking

Electroweak symmetry breaking, explained in section 1.1.1, explains how particles obtain mass through the Higgs mechanism. Yet the origin behind the breaking of the EW symmetry is unclear: the Standard Model does not explain why the Higgs mass parameter μ^2 would be negative. In supersymmetric theories, a positive Higgs mass defined at an energy above the SUSY breaking scale can become negative after running it down to the electroweak scale using RGEs⁹, thus giving the origin of EWSB. It should be noted that this issue is then replaced by the question of what the origin of SUSY breaking is.

Other issues of the Standard Model

The measurement of the anomalous magnetic moment of the muon at BNL, $a_\mu^{exp} = (g - 2)/2 = 11659208.9(5.4)(3.3) \times 10^{-10}$ [53], lies 3.6 standard deviations above the predicted Standard Model value of $a_\mu^{theory} = 11659180.2(0.2)(4.2)(2.6) \times 10^{10}$ [19]. This could be solved in a theory with additional particles in the loops which contribute to a_μ . SUSY would give possible candidates for these particles.

Apart from the aforementioned issues, there are several more outstanding problems with the Standard Model which are not (directly) addressed through the introduction of SUSY. The most important of these are listed below for completeness:

- It does not describe gravity. Attempts to reconcile gravity with the Standard Model, by for instance string theory, have not yet led to a conclusive, working theory. Many string theories contain supersymmetry.
- Although the Standard Model postulates neutrinos to be massless, experiments have shown them to have a finite, small mass [23]. The neutrino flavour oscillations observed cannot be reconciled with the Standard Model either.
- The apparent matter-antimatter asymmetry in the universe. The amount of CP violation allowed in the Standard Model is not enough to explain matter prevailing over antimatter in the universe.

In the next section supersymmetry will be briefly introduced. A full introduction can be found in e.g. [54–56].

1.3 Supersymmetry

Supersymmetry (SUSY) arose as a mathematical possibility in the late 1960s in papers by Miyazawa [57]¹⁰. It was first used in the context of quantum field theories in

⁹The breaking of supersymmetry will be discussed later in this chapter.

¹⁰The name ‘supersymmetry’ stems from Miyazawa’s use of the mathematical concept of supergroups to extend $SU(6)$ to the supergroup $SU(6/21)$, which elements transformed constituent quarks and

the early 1970s, in independent papers by Golfand and Likhtman [59], Volkov and Akulov [60] and Wess and Zumino [61–63]. It quickly rose to fame, with many papers published each year on the subject, due to the ability of supersymmetric theories to solve the issues described in the previous section. This section will first give a short introduction into SUSY, after which a brief view of the theoretical fundamentals is given. Lastly the SUSY phenomenology will be discussed.

Supersymmetry is devised as a symmetry between fermions and bosons within a Lagrangian. Within a supersymmetric theory, for every particle in the Standard Model there exists a partner which carries the same quantum numbers, except that the spin differs by half a unit. In other words, every fermion in the Standard Model has as a corresponding *superpartner* a boson, with the same electrical charge, colour and mass, and vice versa for the Standard Model bosons. It is evident from looking at figure 1.1 that the Standard Model content does not contain fermion-boson superpairs, as there are no fermions and bosons with the same quantum numbers. Therefore supersymmetry requires additional particles. Moreover, the required addition of a second Higgs doublet leads to five Higgs bosons, as will be discussed in section 1.3.2, ending up with double the number of particles in the Standard Model plus an additional four Higgs bosons and four fermionic Higgsinos.

The introduction of the fermion-boson symmetry does come at a price: baryon and lepton number conservation does not hold automatically. Violation of baryon number would lead to issues such as proton decay, which has never been observed – the mean proton lifetime is larger than $10^{31} - 10^{33}$ years [19]. To solve this issue, a new conservation law can be applied, in which *R-parity* is conserved. This multiplicative symmetry is defined using the lepton number L , baryon number B and spin S of a particle:

$$R \equiv (-1)^{3B+L+2S}. \quad (1.31)$$

This equation can be simplified by realising that it means that all Standard Model particles have $R = +1$, while all supersymmetric partners have $R = -1$. The multiplicative behaviour of R-parity leads to an interesting feature when it is conserved: every interaction should have either an even number of supersymmetric particles, or none. This means that heavy supersymmetric particles will always decay into at least one lighter SUSY particle, leading ultimately to a lightest supersymmetric particle which can no longer decay further. This stable lightest supersymmetric particle (LSP), if it is neutral, can thus be a candidate for being (one of) the dark matter particle(s). From here onwards, we consider R-parity to be conserved, although this is not necessarily the case in SUSY: R-parity violating interactions will be baryon or lepton number violating, while for proton decay both the baryon and lepton number need to be violated. Thus by requiring either baryon number or lepton number conservation explicitly, proton decay can be prohibited even when R-parity is not conserved¹¹.

di-quarks into each other. This led to a general definition of $SU(m/n)$ superalgebras, describing a symmetry between m bosons and n fermions: a supersymmetry [58].

¹¹The introduction of explicit B and L conservation could even lead to massive neutrinos [64, 65].

1.3.1 SUSY fundamentals

Supersymmetry in quantum field theories was born out of curiosity: do our current quantum field theories realise all possible symmetries? The symmetry operators incorporated in the Standard Model (e.g. the charge Q) are all Lorentz scalars, i.e. they do not alter the spin of a particle. The Coleman-Mandula theorem [66] states that no conserved operators, or charges, can exist other than Lorentz scalars, the 4-momentum vector operator P_μ which generates space-time translations and the tensor operator $M_{\mu\nu}$ which generate rotations and boosts [55]. Yet there exists a loophole of this theorem: it assumes that all symmetries are bosonic, with conserved charges carrying integer spin. Operators which transform under Lorentz transformations as spinors, i.e. which are fermionic, are exempt from the argument [67]. Such spinor operators Q_a , with a the spinor index, change the spin of a particle:

$$Q_a|J\rangle = |J \pm \frac{1}{2}\rangle. \quad (1.32)$$

Supersymmetry integrates these charges into the Standard model, creating a theory in which the symmetries include transformations that mix bosonic and fermionic states. The supersymmetric algebra that can be constructed from these spinor operators consists of both commutation and anticommutation relations [59, 67]. The simplest form of this supersymmetric algebra is given by [68]:

$$\{Q_a, \bar{Q}_b\} = 2(\gamma^\mu)_{ab}P_\mu \quad (1.33)$$

$$\{Q_a, Q_b\} = 0 \quad (1.34)$$

$$\{\bar{Q}_a, \bar{Q}_b\} = 0, \quad (1.35)$$

with $\bar{Q}_{a,b} = Q_{a,b}^* \gamma^0$ the Dirac conjugate of $Q_{a,b}$. In general there can be N supersymmetric transformation generators in the theory, Q_a^A , where $A = 1, 2, \dots, N$. Yet $N = 1$ SUSY, containing just one set of SUSY transformation generators, is the only case which incorporates chiral representations needed in the Standard Model. SUSY other than $N = 1$ will therefore not be considered.

Equation 1.33 has a tantalising consequence: it says that performing two SUSY transformations after each other is equivalent to the energy-momentum operator. In other words, Q_a can be seen as taking the square root of a space-time translation operator. To be able to do this, space-time itself needs to be extended with extra fermionic degrees of freedom, which are acted on by the SUSY operators. The connection between the normal and fermionic space-time coordinates is achieved by transformations generated by the operators Q . This extended space time is called *superspace*, labelled by coordinates $x_\mu, \theta, \bar{\theta}$. Here θ and $\bar{\theta}$ are anticommuting complex 2-component spinors. The spinor and normal space-time coordinates are connected via SUSY transformations generated by operators Q_a . This interesting concept of the enlargement of normal space-time into superspace will not be covered further in this thesis; see e.g. [55] for more information.

In an $N = 1$ supersymmetric extension of the Standard Model, fermions and bosons

Name	Sparticle fields	Mass eigenstates
Squarks	$(\tilde{u}_L \ \tilde{d}_L)$	$\tilde{u}_1, \tilde{u}_2$
	$\tilde{u}_R$	$\tilde{d}_1, \tilde{d}_2$
	$\tilde{d}_R$	
Sleptons	$(\tilde{\nu}_r \ \tilde{e}_L)$	$\tilde{\nu}_e$
	$\tilde{e}_R$	$\tilde{e}_1, \tilde{e}_2$
Gluino	$\tilde{g}$	$\tilde{g}$
Higgsinos	$(\tilde{H}_u^+ \ \tilde{H}_u^0)$	$\tilde{\chi}_{1,2}^\pm$
	$(\tilde{H}_d^0 \ \tilde{H}_d^-)$	
Wino	$\tilde{W}^\pm, \tilde{W}^0$	$\tilde{\chi}_{1,2,3,4}^0$
Bino	$\tilde{B}^0$	

Table 1.1: *The fields and mass eigenstates of the supersymmetric partners of the Standard Model particles.*

are placed in *supermultiplets*. Complex scalar fields (ϕ) are placed together with two-component chiral fermions (χ) in *chiral supermultiplets*. Spin 1 vector bosons are placed in *gauge supermultiplets* together with their superpartners, spin 1/2 chiral fermions. Thus partners in supermultiplets always differ half a unit in spin.

The fields placed in supermultiplets are related to each other via the SUSY operators Q and $\bar{Q}$: acting with a combination of these operators on one member of the supermultiplet results in the other member, up to space-time translations and rotations. Following the reasoning in ref. [54], these commute with the square mass operator $-P^2$, which itself commutes also with spacetime translations and rotations. Therefore, fields in a supermultiplet should have the same eigenvalue of $-P^2$, and thus have equal mass! Even more, Q and $\bar{Q}$ commute as well with the gauge transformation generators: a fermion-boson pair in a supermultiplet should also have equal charge, weak isospin and colour.

The exact form of SUSY transformations, leading to the SUSY Lagrangian with mass terms, gauge and Yukawa couplings via the superpotential W is outside the scope of this thesis. See for instance [54–56, 69] for detailed discussions. From here onwards we will discuss the particle content of the minimal supersymmetric extension to the Standard Model, and its phenomenology.

1.3.2 Particle content of the MSSM

The particle (or *sparticle*) content introduced in this section, together with the supersymmetric interactions, makes up the minimal supersymmetric Standard Model (MSSM). Although additional supermultiplets can be added to the theory, the MSSM is the minimal extension needed to make the Standard Model supersymmetric. Table 1.1 lists the fields and mass eigenstates of the sparticles in the MSSM.

It can be shown that only the chiral supermultiplets can contain fermions of which the left- and right-handed parts transform differently under gauge transformations [54] – the fermions of the Standard Model, which all possess this feature, must be the fermionic members of the chiral supermultiplets. The scalar partners of the fermions are indicated with a prefix 's': a quark is partnered with a scalar particle called *squark*, short for 'scalar quark' and denoted by $\tilde{q}_L, \tilde{q}_R$, with $q = u, d, s, c, b, t$. Note that the subscript does not denote the helicity of the particles themselves, but of the particles they partner. The partners of the top and to a less degree the bottom quark will be used often: these are called top squark and bottom squark, or *stop* ($\tilde{t}$) and *sbottom* ($\tilde{b}$). Leptons are partnered with *sleptons* ($\tilde{l}$): the partners of electrons, muons and taus are called *selectrons* ($\tilde{e}_L, \tilde{e}_R$), *smuons* ($\tilde{\mu}_L, \tilde{\mu}_R$) and *staus* ($\tilde{\tau}_L, \tilde{\tau}_R$). The left-handedness of neutrinos (neglecting the evidence for massive neutrinos) gives just one scalar partner (*sneutrino*) for each flavour: $\tilde{\nu}_e, \tilde{\nu}_\mu$ and $\tilde{\nu}_\tau$.

After SUSY breaking, discussed in section 1.3.3, trilinear mixing (via the A terms that will be introduced in equation 1.38) allows for mixing of the scalar partners of the left- and right-handed fermions, forming mass eigenstates which are a mixture of both. Since this mixing is proportional to the fermion mass [56], it is largest for sbottoms, stops and staus, while the mixing for lighter squarks and sleptons is negligible. The mixed third generation fermions are denoted by $\tilde{f}_{1,2}$, with $f = t, b, \tau$ for stops, sbottoms and staus, respectively. The lightest state is always given by $\tilde{f}_1$. If the mixing is large $\tilde{t}_1$ will become the lightest of the squarks.

The Standard Model gauge bosons must be the bosonic members of the gauge supermultiplets, as they are spin 1 vector bosons as required. The fermionic partners of the Standard Model gauge bosons are called *gauginos*, where the partner of the gluon is called the *gluino* ($\tilde{g}$), the partners of the $SU(2)_L$ gauge bosons W_i are the *winos* ($\tilde{W}_i$), and the partner of the $U(1)$ gauge field B is the *bino* ($\tilde{B}$). As can be seen, a spin 1/2 superpartner of a boson has 'ino' appended to the Standard Model name.

The Higgs sector is somewhat more complicated. As it has spin 0, the Higgs scalar must occupy a chiral supermultiplet, and its fermionic partner (a *Higgsino*) will contribute to gauge anomalies. This is a problem: in the Standard Model these anomalies cancel exactly, yet they are reintroduced by inserting the fermionic Higgs partner. This can be solved by adding a second Higgs supermultiplet with opposite quantum numbers: the contributions of the two Higgsinos to the gauge anomalies will thus cancel. Furthermore, unlike in the Standard Model it is not possible in a supersymmetric theory to give both the up- and down-type quarks mass with the same Higgs doublet. In the Standard Model the charge conjugate of the Higgs doublet generates mass for the up-type quarks, however charge conjugates are not allowed in SUSY. Introducing the second Higgs supermultiplet solves this issue: one scalar doublet couples to up-type quarks, which we will call $H_u = (H_u^+, H_u^0)$, while the other ($H_d = (H_d^0, H_d^-)$) couples to down-type quarks and charged leptons. This sums up to four complex scalar fields, or eight real degrees of freedom. Just like in the Standard Model, three of these scalar fields will be used via EWSB to generate the massive Z

and $W^\pm$ bosons, leaving *five* real scalar Higgs mass eigenstates: two CP-even neutral scalars h and H , one CP-odd neutral scalar A and two charged scalars $H^\pm$. The vacuum expectation values v_u and v_d of the two neutral components of the Higgs doublets determine together the mass of the W boson [56]:

$$M_W^2 = \frac{g^2}{2}(v_u^2 + v_d^2). \quad (1.36)$$

In other words, $(v_u^2 + v_d^2) = v^2$. Since v^2 is fixed, there are only 3 free parameters left to determine the total Higgs sector. Two of these are usually taken to be $\tan \beta = \frac{v_u}{v_d}$ and the mass of the CP-odd Higgs boson A , m_A . It can be shown that the value of μ has been eliminated [55], yet the sign of μ has not. This is the final parameter characterising the Higgs sector.

The superpartners of the Higgs doublets are the Higgsino doublets $\tilde{H}_u = (\tilde{H}_u^+, \tilde{H}_u^0)$ and $\tilde{H}_d = (\tilde{H}_d^+, \tilde{H}_d^0)$, adding up to four fermionic Higgsino fields. Similar to the mixing which occurs in the Standard Model between the gauge bosons leading to the $W^\pm$, γ and Z bosons, in the MSSM the gauginos and Higgsinos mix into four neutral and four charged states. As the gluino is a colour octet, it does not participate in the mixing. When the neutral scalar Higgs fields (H_u^0 and H_d^0) acquire a non-zero vacuum expectation value, mixing will occur and combinations of the neutral Higgsinos ($\tilde{H}_u^0$ and $\tilde{H}_d^0$) with the neutral electroweak gauginos ($\tilde{B}$ and $\tilde{W}_3$) will be generated. These new neutral mass eigenstates are the *neutralinos* and denoted by $\tilde{\chi}_i^0$ with $i = 1, 2, 3, 4$, where the index runs from lightest to heaviest particle. The same occurs for the charged Higgsinos ($\tilde{H}_u^+$ and $\tilde{H}_d^-$) which mix with the charged electroweak gauginos ($\tilde{W}^\pm$) to form *charginos*: $\tilde{\chi}_i^\pm$ with $i = 1, 2$. The lightest neutralino is usually assumed to be the lightest supersymmetric particle (LSP), assuming R-parity is conserved and there is no lighter gravitino. As it meets all requirements, it makes for a good dark matter candidate [54, 70].

1.3.3 Supersymmetry breaking

As discussed in the previous sections, if supersymmetry would be an exact symmetry, superpartners would have equal mass. Yet no superpartner of any of the Standard Model particles has been discovered at low energies, thus SUSY must be a broken symmetry, with a mechanism for breaking the symmetry at some higher energy scale. The symmetry is expected to be broken spontaneously: with a Lagrangian invariant under SUSY transformations while the vacuum state of the symmetry is not invariant [54]. The mechanism behind such a spontaneous symmetry breaking is not known – many models have been proposed, such as gravity mediated SUSY breaking [71], gauge mediated SUSY breaking (GMSB) [72] and anomaly mediated SUSY breaking (AMSB) [73]. As there is no consensus over which model should be realised in nature, the MSSM takes the practical way forward and introduces the SUSY breaking explicitly: the MSSM is assumed to be an effective low energy theory, and terms which break the symmetry are added explicitly to the MSSM Lagrangian. The precise origin of the broken symmetry is left as an open question.

The hierarchy problem provides extra input for these breaking terms. The Higgs mass corrections given SUSY are of the form

$$\delta m_h^2 \sim (\lambda_S - |\lambda_f|^2) \Lambda^2 + \dots, \quad (1.37)$$

where λ_S (λ_f) are the couplings between the Higgs and the SUSY scalar and between the Higgs and the Standard Model fermion, respectively. As we do not want to reintroduce the quadratic divergences which caused the hierarchy problem, the relationship between the dimensionless couplings (λ_S and λ_f) should be unchanged through SUSY breaking. If SUSY breaking interactions alter either λ_S or λ_f , the mass shift in equation 1.37 becomes large again. It can be shown that this means we need ‘soft’ SUSY breaking, where the SUSY Lagrangian is supplemented with supersymmetry violating terms which have mass dimensions between 1 and 3. These possible soft operators can thus only include mass terms, bilinear mixing terms and trilinear scalar mixing terms, and the Lagrangian describing the soft SUSY breaking for the first generation is given by [55]:

$$\mathcal{L}_{soft} = -\frac{1}{2}(M_3\tilde{g}\tilde{g} + M_2\tilde{W}\tilde{W} + M_1\tilde{B}\tilde{B} + \text{h.c.}) \quad (1.38)$$

$$- m_{\tilde{Q}_1}^2 \tilde{Q}_1^* \tilde{Q}_1 - m_{\tilde{u}}^2 \tilde{u}_R^* \tilde{u}_R - m_{\tilde{d}}^2 \tilde{d}_R^* \tilde{d}_R + m_{\tilde{L}_1}^2 \tilde{L}_1^* \tilde{L}_1 + m_{\tilde{e}}^2 \tilde{e}_R^* \tilde{e}_R \quad (1.39)$$

$$- m_{H_u}^2 |H_u|^2 + m_{H_d}^2 |H_d|^2 - (b H_u H_d + \text{h.c.}) \quad (1.40)$$

$$- (a_u \tilde{d}_R^* \tilde{Q}_1 H_u - a_d \tilde{u}_R^* \tilde{Q}_1 H_d - a_e \tilde{e}_R^* \tilde{L}_1 H_d + \text{h.c.}). \quad (1.41)$$

Note that in the full expression for all generations the squark and slepton masses and trilinear couplings are matrices rather than scalar numbers. The first line has the gaugino masses for each gauge group (M_1 , M_2 , M_3); the second line contains the squark and slepton mass terms; the third line the Higgs mass terms; while the last line contains the triple scalar couplings. Note that $\tilde{Q}_1$ and $\tilde{L}_1$ denote the first generation left-handed squark and slepton doublet, respectively, as shown in table 1.1. The explicit mass terms for the gauginos and scalars effectively break the mass degeneracy between them and their Standard Model partners.

1.3.4 SUSY Phenomenology

The addition of $\mathcal{L}_{soft}$ has had some unwanted consequences: unlike the unbroken theory with just one free parameter (μ) added to the Standard Model, we need to add 105 additional free parameters to manage all the new masses and couplings. Even though many of these parameters have constraints from measurements on CP violation, flavour changing neutral currents and other precision measurements, there is a large SUSY landscape with a plethora of possibilities for the masses, couplings and decay modes of the SUSY particles. Searches for SUSY are thus faced with a difficult task: instead of searching for a single signature of a theory, many different signatures need to be evaluated. The signature in the ATLAS detector depends vastly on the masses of the SUSY particles, and the SUSY breaking mechanism realised in nature

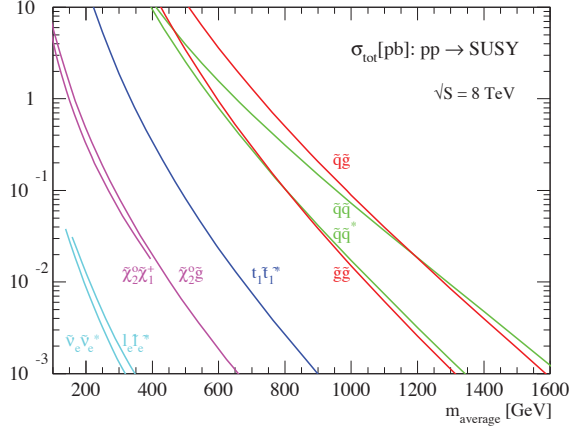


Figure 1.7: The sparticle pair production cross section as a function of their mass at the LHC with $\sqrt{s} = 8$ TeV, in picobarn [74].

(assuming SUSY is realised at all). In ATLAS the search has been mostly subdivided by reconstructed objects observed in the detector. This thesis focusses mainly on the search for SUSY with jets and missing transverse momentum, where no leptons are present. The SUSY phenomenology discussed below will therefore focus on these signatures. The models which are used in chapters 3 and 5 to interpret the results will be shortly introduced here, although these are by no means all available models.

The most abundantly produced new particles will be those which couple strongest to the coloured particles in the proton (unless the new particles are too heavy to be produced). In SUSY, these are the gluinos and squarks, as can be seen from the production cross sections, shown in figure 1.7 for the production of SUSY particles at the LHC running at $\sqrt{s} = 8$ TeV. These are mostly produced through interactions with gluons or valence quarks, as the production of high mass squarks and gluinos requires interacting partons with high momentum fraction of the incoming protons. Requiring R-parity conservation and a neutral particle as the LSP, such as the lightest neutralino, has some implications for the decays of these particles. If $m_{\tilde{q}} \ll m_{\tilde{g}}$, squark-pair production dominates, with dominant decay of either squark to a quark plus a LSP, $\tilde{q} \rightarrow q\tilde{\chi}_1^0$, although decay to a chargino and quark can occur as well, with subsequent decay into a lepton. For $m_{\tilde{g}} \ll m_{\tilde{q}}$, gluino production dominates, where the tree level decay of the gluino can be into a quark-antiquark pair plus neutralino ($\tilde{g} \rightarrow q\bar{q}\tilde{\chi}_1^0$) or a squark and a quark, with subsequent squark decay ($\tilde{g} \rightarrow \tilde{q}^*q \rightarrow \bar{q}q\tilde{\chi}$). The one-loop decay of gluinos into a gluon and neutralino can occur as well but is suppressed.

When the produced particles are heavy, especially with large mass differences between the produced squark/gluino and the SUSY particle in the decay, the quarks produced in $\tilde{g}$ or $\tilde{q}$ decay will be heavily boosted, leading to high momentum jets in the detector (see chapter 2). The SUSY particle in that decay can be the LSP, or another SUSY particle such as a chargino which subsequently decays, leading eventually

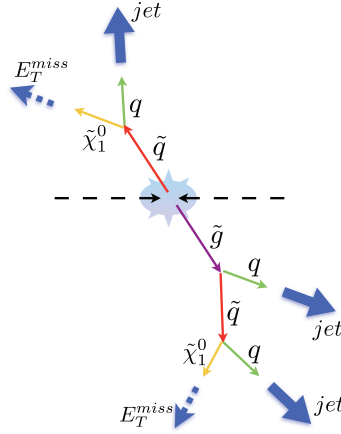


Figure 1.8: Illustration of the production of a squark and gluino, and their decay into quarks and neutralinos, leading to jets and $\cancel{E}_T$.

to the stable LSP which escapes the detector. The LSP thus leads to missing transverse energy. Squark-pair production, squark-gluino production and gluino-pair production produces a signature of two, three and four high momentum jets, respectively, with missing transverse energy, and possible other decay products such as leptons, coming from the potential subsequent decays. The production and direct decay of a squark and a gluino is illustrated in figure 1.8.

Minimal supergravity and CMSSM

The various SUSY breaking mechanisms all have specific consequences for SUSY phenomenology. Historically, minimal supergravity (mSUGRA) [75, 76] and the similar constrained MSSM (CMSSM) [75, 77] have had the most attention. In supergravity models, gravity is the mediating force between the hidden, unbroken sector and the MSSM, generating the soft breaking terms naturally. In the minimal theory, several simplifying assumptions are made on top of the full supergravity Lagrangian. Not only are the couplings unified at the GUT scale, also mass unification occurs. At the GUT scale all gauginos have a common mass $m_{1/2}$, where $M_3 = M_2 = M_1 = m_{1/2}$, with M_1, M_2, M_3 as in equation 1.38. Likewise, all scalars obtain a common mass m_0 : $m_{\tilde{q}}^2 = m_{\tilde{l}}^2 = m_{H_u}^2 = m_{H_d}^2 = m_0^2$. A third simplification is that the trilinear couplings a_u , a_d and a_e (and their second and third generation counterparts) in equation 1.41 are assumed to be equal, and given by a common coupling A_0 at the GUT scale. With these assumptions the full MSSM spectrum can be predicted, using the RGEs, in terms of just five parameters: the two common masses m_0 and $m_{1/2}$; the common trilinear coupling A_0 ; the ratio of the two Higgs vacuum expectation values, $\tan \beta$; and the sign of μ , the Higgs mass parameter. Note that the common masses for gauginos and scalars only hold at the GUT scale – after SUSY breaking, the masses and couplings evolve when the energy scale is run down to the electroweak scale with the RGEs,

resulting in non-degenerate masses.

The dependency of the mass spectrum of mSUGRA/CMSSM on the five parameters is varied. As a rule of thumb, all masses increase with increasing $m_{1/2}$, while only squark and slepton masses depend on m_0 . Increasing $\tan\beta$ has the effect of larger mixing in the third generation sfermions, leading to mass differences between $\tilde{t}_1$ and $\tilde{t}_2$, and between $\tilde{\tau}_1$ and $\tilde{\tau}_2$ [78]. A_0 has a small effect on the lightest neutral Higgs mass, while for increasing $|A_0|$ mass of the lightest stop, $\tilde{t}_1$, decreases due to mixing. Furthermore, in mSUGRA the masses of the gluino, lightest chargino and lightest neutralino are related from GUT scale unification and the RGEs, with an approximate relation of $m_{\tilde{g}} : m_{\tilde{\chi}_1^\pm} : m_{\tilde{\chi}_1^0} \sim 6 : 2 : 1$ [54].

Simplified models

The complexity of SUSY and its specific models is a disadvantage when interpreting results of an analysis. A certain lower limit on the gluino mass in one model might be a much looser limit in a different model. Instead of interpreting the results of an analysis in a specific model such as the CMSSM, one can also try a more general approach. Both ATLAS and CMS also set limits in a *simplified model* framework [79, 80], in which more phenomenological models are used.

A simplified model is based on an effective Lagrangian, which only describes a small set of particles and interactions, production cross sections and branching ratios. All other particles are integrated out of the theory, or in other words, are set to such a high mass that they do not participate in interactions. In this sense, simplified models are not physically sensible models, nor are they completely model-independent. However, they come with some benefits.

Firstly, interpreting an analysis in terms of just one type of process, where the mass parameters of the participating particles vary, helps to identify the sensitivity of the search. The kinematics of a process depends heavily on the masses and mass differences of the participating particles, and thus the reconstruction efficiencies do as well. By analysing simple topologies, kinematic ranges and topologies can be identified where the analysis can be improved.

Secondly, although the model itself may not be physically 'sensible', limits on several simplified models can be combined to set a constraint on a variety of 'real' models. This is done as follows: for each of the topologies i which are available in the general model, a simplified model should be found, with its corresponding production fraction P_i : the fraction of events in the general model of topology i . The number of expected events N of the general model in a given signal region can then be found as a product of the total cross section, the integrated luminosity and the sum of the experimental efficiencies ϵ_i , production fraction and branching ratios BR_i for topology i – or in an equation [81]:

$$N = \sigma_{tot} \times \mathcal{L}_{int} \times \sum_{i \in \text{Simp. Mod.}} \epsilon_i \times P_i \times BR_i^2. \quad (1.42)$$

Finally, simplified models can be used to determine quantum numbers of new

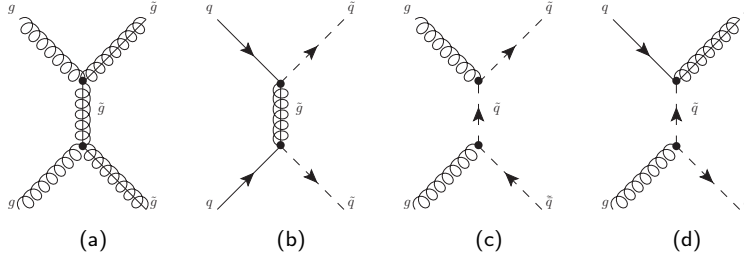


Figure 1.9: Tree level Feynman diagrams of (a) gluino-pair production, (b) squark-pair production, (c) squark-antisquark production and (d) associated squark-gluino production. Apart from the t -channel diagrams shown, there are also s -channel diagrams for all but $\tilde{q}\tilde{q}$ production [83].

physics, when discovered [82]. In this case, it is useful to distinguish features of the experimental data with as few parameters as possible, thus constraining masses and cross sections without external assumptions.

Models which are used in the search for squarks and gluinos in inclusive hadronic searches in ATLAS have either a pair of squarks, gluinos or both being produced directly at the collision, while no other sparticle is allowed to be produced. Tree level Feynman diagrams of the t -channel production of these particles are shown in figure 1.9. The most straightforward models have these sparticles decay directly to quarks and an LSP, while it can also occur with one or more intermediate charginos. To be precise, the four simplified models, proposed in ref. [79], which are used in the analysis described in chapter 5, are:

- Pair production of gluinos, which each decay directly to two quarks and one LSP, $\tilde{g} \rightarrow q\bar{q}\tilde{\chi}_1^0$. See figures 1.9 (a) and 1.10 (a);
- Production of a squark-antisquark pair, which each decay directly to one quark and one LSP, $\tilde{q} \rightarrow q\tilde{\chi}_1^0$. See figures 1.9 (c) and 1.10 (b);
- Pair production of gluinos, which each decay directly to two quarks and a chargino. The chargino subsequently decays into a W boson and an LSP: $\tilde{g} \rightarrow q\bar{q}\tilde{\chi}_1^\pm \rightarrow q\bar{q}W^\pm\tilde{\chi}_1^0$. See figures 1.9 (a) and 1.10 (c);
- Production of a squark-antisquark pair, which each decay directly to one quark and a chargino. The chargino subsequently decays into a W boson and an LSP: $\tilde{q} \rightarrow q\tilde{\chi}_1^\pm \rightarrow qW\tilde{\chi}_1^0$. See figures 1.9 (c) and 1.10 (d).

The absence of pure squark-pair production can be understood from the Feynman diagram in figure 1.9 (b). Without a possible s -channel diagram, the process is always mediated by a virtual gluino. However, to study pure squark production this gluino is decoupled from the theory, by setting it to an infinite mass. Therefore this diagram cannot contribute, removing possible squark-pair production.

In ATLAS, these simplified models are all defined in a two dimensional plane: events are simulated for points in the $m_{\tilde{q}}-m_{LSP}$ or $m_{\tilde{g}}-m_{LSP}$ mass planes. In the case of

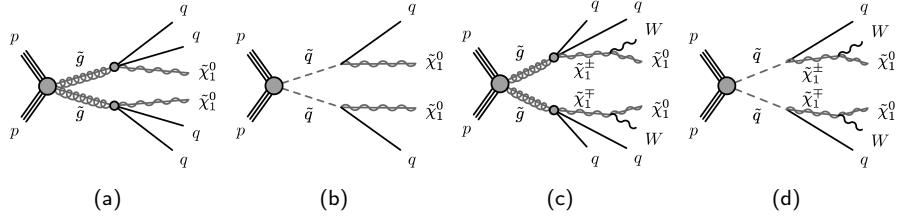


Figure 1.10: Illustrations of simplified models for the decays of (a) $\tilde{g}\tilde{g}$ production, both decaying as $\tilde{g} \rightarrow q\bar{q}\tilde{\chi}_1^0$; (b) $\tilde{q}\tilde{q}$ production, both decaying as $\tilde{q} \rightarrow q\tilde{\chi}_1^0$; (c) $\tilde{g}\tilde{g}$ production, both decaying as $\tilde{g} \rightarrow q\bar{q}\tilde{\chi}_1^\pm \rightarrow q\bar{q}W^\pm\tilde{\chi}_1^0$; (d) $\tilde{q}\tilde{q}$ production, both decaying as $\tilde{q} \rightarrow q\tilde{\chi}_1^\pm \rightarrow qW^\pm\tilde{\chi}_1^0$.

the models with an intermediate chargino, the mass of the chargino is assumed to be halfway between the squark/gluino and the LSP mass. In these models the kinematics depend significantly on the mass difference between the squark or gluino and the LSP. Smaller mass differences, so-called *compressed spectra*, generally lead to jets with lower transverse momentum and lower missing transverse energy. In chapter 5 we will see that this affects exclusion limits in these models.

A step closer to a physical scenario is to add several of these processes in one model. In a so-called phenomenological MSSM model of squark/gluino production, all four processes shown in figure 1.9 can be produced: $\tilde{q}\tilde{q}$, $\tilde{g}\tilde{g}$, $\tilde{q}\tilde{q}$ and $\tilde{g}\tilde{g}$. All of these are required to decay directly to quarks and neutralinos. Since there are three participating particles, the parameter space is three-dimensional. To be able to interpret the model in a clear way, two-dimensional slices in the squark-gluino mass plane are made for several values of the LSP mass.

1.3.5 Constraints on SUSY

After having been introduced to possible SUSY signatures in the previous section, we turn our attention to the constraints which have already been set on SUSY, with focus on R-parity conserving SUSY. Many different experiments have already searched for signs of SUSY, directly or indirectly; amongst others colliders (such as LEP and the TEVATRON), precision measurements and dark matter detection experiments. As there are many different SUSY breaking mechanisms and therefore SUSY implementations, many searches have concentrated on specific SUSY scenarios, with CMSSM being the most popular model. Constraints set using this specific model are not model-independent, making the specific mass constraints listed below susceptible to change when considered in a different model.

Previous collider constraints

Before the LHC started¹², the most stringent direct limits on SUSY were set by LEP and the TEVATRON. CDF set a lower limit at 95% confidence in the CMSSM of 392 GeV on the mass of squarks and gluinos when they are degenerate in mass [84], while the D0 limit is set at 379 GeV and 308 GeV on the mass of squarks and gluinos, respectively [85]. The strongest limits on gauginos were set by a LEP combination, with a lower limit on the mass of the lightest neutralino and chargino (in the CMSSM) of 47 GeV and 92 GeV, respectively [86,87]. The mass limit on charginos is 103.5 GeV for models with high mass difference between $\chi^\pm$ and $\tilde{\chi}_1^0$, and high electron sneutrino mass. However, without the assumption of CMSSM, no limit has been set on $\tilde{\chi}_1^0$.

Precision measurements

Precision measurements of electroweak and flavour physics contribute to SUSY constraints as well. The lack of observation of flavour changing neutral currents (FCNCs) indicates that the soft mass terms in SUSY need to be diagonal, which is intrinsically the case for gauge- and anomaly-mediated SUSY breaking mechanisms. The measurement of the rare decay $B_s \rightarrow \mu\mu$ by LHCb, with a branching ratio of $BR(B_s^0 \rightarrow \mu^+\mu^-) = (2.9_{-1.0}^{+1.1}) \times 10^{-9}$ [88], which is in agreement with the Standard Model prediction of $(3.23 \pm 0.27) \times 10^{-9}$ [89], lead to initial speculation that it would be a blow to SUSY. However, although high $\tan\beta$ models may have large enhancements to the Standard Model prediction of that branching ratio, the constraint coming from the $B_s \rightarrow \mu\mu$ measurement is strongly dependent on other MSSM parameters such as m_A and μ , and most SUSY scenarios are compatible with the result [90].

Furthermore, the previously mentioned measurement of the anomalous magnetic moment of the muon a_μ at BNL lies 3.6 standard deviations above the predicted Standard Model value [19]. This discrepancy can be explained through a SUSY contribution to a_μ , if the SUSY model has either light smuons or large $\tan\beta$ [91,92]. Additionally, positive values of μ are favoured.

Astrophysical constraints

Astrophysical observations mainly provide constraints on the LSP as a weakly interacting massive particle (WIMP) candidate, in other words a dark matter candidate. The requirement of an electrically neutral dark matter candidate means SUSY scenarios with the lightest stau as the LSP cannot provide a WIMP candidate, and are therefore less physically relevant. Assuming dark matter particles are a thermal relic from the early universe, the current dark matter abundance can be explained by a thermal ‘freeze out’: whereas in the early universe dark matter particles annihilated abundantly, at a certain moment in time the expansion of the universe caused a balance in the annihilation. The current relic density measured by WMAP, $\Omega_\chi h^2 = 0.1148 \pm 0.0019$ [49], and PLANCK, $\Omega_\chi h^2 = 0.1199 \pm 0.0027$ [50], requires a dark matter annihilation cross

¹²As searches for SUSY in ATLAS are the topic of this thesis, constraints from either ATLAS or CMS on SUSY are ignored here.

section higher than what is possible in large parts of the SUSY phase space. However, it can be achieved in various regions of parameter space [93].

Many experiments searching for WIMP annihilation or nuclear recoils with WIMPs have been conducted in recent years. Although several have found hints of signals for light (< 50 GeV) WIMP candidates, such as annual modulation in the signal from DAMA/LIBRA [94] and CoGENT [95], a possible WIMP signal at CRESST-II [96], they contradict one another when interpreted in WIMP mass. Furthermore, constraints from the XENON100 [97] and LUX [98] experiments on the WIMP-nucleon cross section contradict all of the above.

Constraints from the Higgs mass

The discovery of a Higgs boson with a mass of ~ 126 GeV [22] has additional consequences for the SUSY landscape. The tree-level lightest Higgs mass m_h has an upper bound in the MSSM of $m_h \leq m_Z |\cos(2\beta)| \sim 90$ GeV. However, radiative corrections can have large contributions to m_h and can increase the lightest Higgs mass to $m_h \sim 130$ GeV, with the dominant effects coming from loops involving third generation quarks and squarks, which are proportional to the relevant Yukawa couplings [99, 100]. The approximate one-loop equation is given by:

$$m_h^2 = m_Z^2 \cos^2(2\beta) + \frac{3m_t^4}{4\pi^2 v^2} \left\{ \log \left(\frac{M_S^2}{m_t^2} \right) + \frac{X_t^2}{M_S^2} \left(1 - \frac{X_t^2}{12M_S^2} \right) \right\}, \quad (1.43)$$

with m_t the top mass, and M_S the mass of the SUSY particles $M_S = \sqrt{m_{\tilde{t}_1} \cdot m_{\tilde{t}_2}}$. The left-right stop mixing parameter X_t is given by $X_t = a_t - \mu \cot \beta$ where a_t is the trilinear coupling for the top sector.

Given the Higgs mass measurement, constraints can be set on $\tan \beta$, X_t and M_S in the MSSM: the ‘no-mixing’ scenario, with $X_t \sim 0$, is ruled out for $M_S \lesssim 3$ TeV, requiring very high stop masses for both the left- and right-handed stop [101]. In general, large values of X_t/M_S are required: for intermediate mixing, $1 \text{ TeV} \lesssim M_S \lesssim 3 \text{ TeV}$ together with moderately large values of $\tan \beta \gtrsim 20$ are needed [102], while most models with near maximal mixing ($A_t/M_s \sim \sqrt{6}$), survive. These will have a large mass splitting between the two stops, where the lightest stop can be as light as 200 GeV [100].

Among the constrained MSSM models, GMSB and AMSB are more heavily affected by the Higgs measurement. They can reach a maximal lightest Higgs mass of 121.5 GeV and 121 GeV, respectively [101, 102], and are therefore ruled out as feasible SUSY models in their current formulation. The CMSSM can just about produce a 125 GeV Higgs, at the cost of high stop masses, high $\tan \beta$ and large fine tuning.

Fits to data in constrained models

Using the above constraints, a fit can be performed to obtain the most likely values of SUSY masses. Prior to the LHC searches, fits to experimental data from measurements on a_μ , direct limits from LEP and cosmological constraints lead to best fit values for

squarks and gluinos between 600 - 800 GeV in the CMSSM, while masses up to 2.5 TeV are still allowed [103]. The best fit of the LSP mass was around 100 GeV before the start of the LHC. These results are largely determined by the results on a_μ and dark matter detection experiments. Note that these fits are not model independent, but are performed in specific SUSY scenarios, in this case the CMSSM. In the final discussion chapter of this thesis (chapter 6) the values after the two years of LHC data are discussed.

The LHC and the ATLAS detector

To study the direct production of new massive particles, an experimental set-up is needed where particles collide with high centre-of-mass energy, and where the final state particles are detected with high precision. To this end, collisions provided by the LHC accelerator are studied by the ATLAS experiment. This chapter describes the configuration and performance of the accelerator in the first section, and that of the ATLAS detector in the succeeding sections.

2.1 The LHC accelerator

The Large Hadron Collider [104] is a circular proton-proton accelerator and collider situated at the CERN laboratory, near the city of Geneva at the border of France and Switzerland. It is placed in the 26.7 km long tunnel constructed in the 1980s for the Large Electron Positron (LEP [105]) collider, which was shut down in the year 2000. The circular tunnel sits approximately 100 m below the surface, at 170 m below the surface at the Jura-side of the ring and 45 m below the surface at the side of Lake Geneva (Lac Léman). The LHC is designed to conduct proton-proton collisions which reach a centre-of-mass energy of 14 TeV with a luminosity of $10^{34} \text{cm}^{-2} \text{s}^{-1}$. Yet due to an accident at the initial start-up in September 2008 it has been decided to run with a lower centre-of-mass energy of $\sqrt{s} = 7$ TeV in 2010-2011, which was increased to $\sqrt{s} = 8$ TeV in 2012. The accelerator will be upgraded to the full design energy only in 2015. This still means that the LHC is the world's highest energy collider: the TEVATRON collider, situated at Fermilab near Chicago, ran up to September 2011 at $\sqrt{s} = 1.96$ TeV, while the Super Proton-antiproton Synchrotron (SPPS) situated at CERN in the 1980s collided protons and anti-protons at $\sqrt{s} = 540$ GeV, which was upgraded to 630 GeV. Finally, the previously mentioned LEP accelerator collided electrons and positrons with centre-of-mass energies up to 209 GeV. Apart from protons the LHC also has a program to accelerate and collide lead nuclei, as well as protons with lead nuclei. These are outside the scope of this thesis.

The protons are produced by stripping off electrons from hydrogen atoms, thus ionising these atoms. Before arriving at the LHC accelerator ring these protons traverse

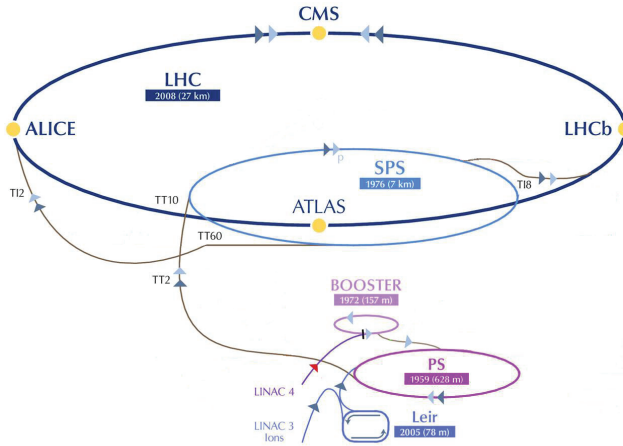


Figure 2.1: *The CERN accelerator complex: the chain of LHC pre-accelerators, starting with the LINAC, ending up in the LHC. The four main LHC experiments are shown in the large ring. Taken from [106].*

a series of pre-accelerators to prepare them for the LHC injection energy of 450 GeV. The first step is acceleration to 50 MeV in the Linear Accelerator (LINAC). From here the protons are accelerated in a series of three circular pre-accelerators: to 1.14 GeV in the Proton Synchrotron Booster (PSB); to 26 GeV in the Proton Synchrotron (PS); and finally the Super Proton Synchrotron (SPS) accelerates the protons to an energy of 450 GeV, with protons travelling with nearly the speed of light. During this sequence the protons are arranged into small packets – the design for these proton bunches is to contain $\sim 10^{11}$ protons each, with a separation of 25 ns between them. The SPS finally injects these bunches into the LHC in both the clockwise and the anti-clockwise direction. A schematic view of the CERN accelerator complex is shown in figure 2.1.

To keep the proton beams in their circular path in the LHC, 1232 superconducting dipole magnets are used, providing a magnetic field strength of 8.33 T. To achieve this they are cooled with liquid helium to 1.9 K. As the existing tunnel is 3.7 m in diameter, there is no room for a separate clockwise and anti-clockwise ring. Both beams thus need to travel through two separate vacuum beam-pipes, and are bent by separate dipole systems, all within one single iron yoke and using one cryostat.

To accelerate the protons, Radio Frequency (RF) cavities are placed at one point in the ring. These accelerate the protons each revolution by use of an oscillating electric field of 2 MV with a frequency of 400 MHz. Accelerating the proton beams from their injection energy of 450 GeV to the final 4 TeV took approximately 15 minutes at the end of 2012.

At certain points in the LHC ring the two proton beams cross each other, enabling proton-proton collisions. Around these points particle detectors are built: there are four large and two smaller detectors to conduct experiments with the LHC data. The two largest are ATLAS [107] and CMS [108], situated at opposite sides of the LHC

	2010	2011	2012	design
$\sqrt{s}$	7 TeV	7 TeV	8 TeV	14 TeV
Delivered integrated luminosity [$\text{fb}^{-1}/\text{year}$]	0.049	5.62	23.25	80-120
ATLAS recorded integrated luminosity [fb^{-1}]	0.045	5.25	21.74	–
Max. number of colliding bunches	348	1854	1380	2808
Bunch spacing [ns]	75	75/50	50	25
Peak instantaneous luminosity [$\text{cm}^{-2}\text{s}^{-1}$]	2.07×10^{32}	3.65×10^{33}	7.73×10^{33}	10^{34}
$\langle \mu \rangle$	–	9.1	20.7	19.2
Peak $\langle \mu \rangle$	3.78	32.21	69.49	–

Table 2.1: *The LHC performance during its three years of data taking. The various rows give the centre-of-mass energy $\sqrt{s}$, the total integrated luminosity per year, both the amount delivered by the LHC to ATLAS and that which was actually recorded, the maximum number of colliding bunches that year, the bunch spacing used, and the maximum instantaneous luminosity. The last two rows give the average number of interactions per bunch crossing $\langle \mu \rangle$ and the maximum number of interactions per bunch crossing (peak μ). The values represented by a dash were not available. The design values are taken from [113].*

ring. Both ATLAS and CMS are so-called 4π general purpose detectors. They are designed to detect the full-range of possibilities from proton-proton collisions, including precision measurements of known SM processes, searches for the SM Higgs boson as well as searches for phenomena which cannot be explained by the Standard Model. On either side between them are two detectors built to study specific phenomena: the ALICE [109] and LHCb [110] detectors. The former is dedicated to the heavy-ion collisions, studying the properties of high energy density nucleus-nucleus interactions. Its goal is to form and observe a new matter phase, the quark-gluon plasma. The LHCb experiment is devoted to B-hadron physics, studying CP-violation and rare decays of the B-hadrons.

The two smallest detectors are TOTEM [111] and LHCf [112]. The former studies the total proton-proton cross section, which measurement is used by all other experiments, while the latter performs measurements of neutral particles emitted in the very forward region of the collisions, i.e. very close to the direction of the beams.

2.1.1 Short description of LHC operations

In 2010 and 2011, the first two full years of operation, the LHC accelerated the proton beams to 3.5 TeV each. This was increased in 2012 to 4 TeV which was decided to be the maximum energy the accelerator could reach while remaining relatively safe from magnet quenches. This will be increased to near its design energy of 7 TeV per beam after a shut-down in the years 2013 and 2014. The instantaneous luminosity,

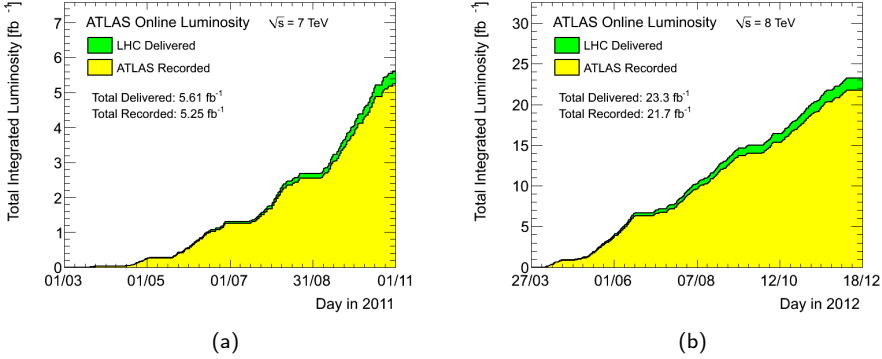


Figure 2.2: The integrated luminosity in 2011 (left) and 2012 (right) as a function of the day in the year. The green histogram shows the data delivered to ATLAS, while the yellow shows the data recorded by ATLAS. Note the different scale used on the y -axis.

proportional to the number of protons per beam and inversely proportional to the area of interaction, has increased dramatically over the three years from $2.07 \times 10^{32} \text{ cm}^{-2}\text{s}^{-1}$ at the end of 2010 to $7.73 \times 10^{33} \text{ cm}^{-2}\text{s}^{-1}$ at the end of 2012. This was mainly due to increasing the number of protons per bunch, increasing the number of bunches in the machine and decreasing the bunch diameter. Although the goal of a 25 ns bunch spacing has not yet been reached, during 2010 and the first months of 2011 the LHC operated with a bunch spacing of 75 ns, while from the second half of 2011 the LHC has operated with a 50 ns bunch spacing. The number of bunches which crossed each other at the ATLAS interaction point increased from 113 at the end of 2010 to 1377 at the height of 2012 data taking.

As the bunch intensity drops during a run due to proton-gas and proton-proton collisions, the LHC runs for $\sim 10 - 20$ hours with one fill, after which it needs to be refilled. Furthermore, time is needed for maintenance. This has a result that the LHC delivers collisions about $\sim 30\%$ of the time. The total delivered integrated luminosity in this time was 49 pb^{-1} in 2010, 5.6 fb^{-1} in 2011, both at $\sqrt{s} = 7 \text{ TeV}$, and 23.3 fb^{-1} in 2012 at $\sqrt{s} = 8 \text{ TeV}$. This performance of the LHC in the first three years of operation is summarised in table 2.1, and the accumulation of data during these years is shown in figure 2.2, with the integrated luminosity delivered to ATLAS as a function of time in green, and recorded by ATLAS in yellow. The left figure shows the 2011 values, and the right figure those of 2012.

2.2 Overview of ATLAS

The ATLAS (A Toroidal LHC ApparatuS) detector [107] is a colossal construction, measuring 44 m in length and 25 m in height, and weighing 7000 tonnes. The size is mostly due to the toroid magnet used for the momentum measurement by the muon

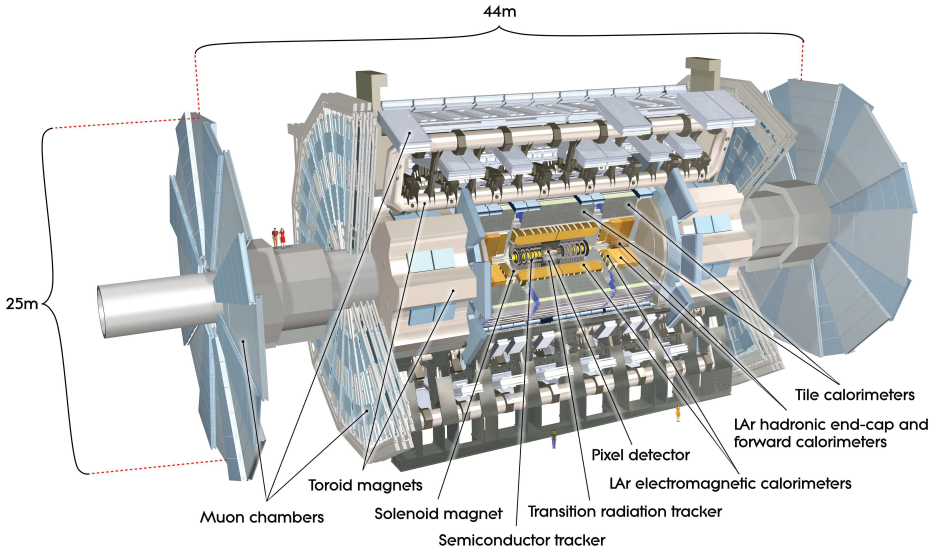


Figure 2.3: Cutaway view on the ATLAS detector. The various detector elements are denoted by different colours and their description. Note the pair of people drawn to scale on the left. Taken from [107].

spectrometer.

The ATLAS detector consists of several sub-detectors, which are placed around the interaction point, as seen from the illustration in figure 2.3. These are organised in a *barrel* region, which consists of concentric cylindrical layers, and two *end-cap* regions where the detector-elements are constructed in wheels which close off the barrel cylinders. The detector can be categorised in three main layers, which all have a specific purpose, and each consist of multiple sub-detectors. The *inner detector* (ID) sits closest to the interaction point, and has as its main purpose the tracking of charged particles and the measurement of their momentum, through their deflection in the magnetic field of a *solenoid magnet*. The ID is discussed further in section 2.4. The electromagnetic and hadronic *calorimeters* absorb hadrons, photons and electrons, and measure their energy. This is discussed in section 2.5. The remaining particles (apart from the weakly interacting neutrino) to fly through the calorimeters are mostly muons, which are tracked by ATLAS's outermost sub-detector, the *muon spectrometer* (MS). It tracks the trajectory of the muons through the magnetic field generated by the *toroid magnet* and thus measures their momentum. The muon spectrometer is discussed further in section 2.6. The magnet system providing the necessary bending power is described in section 2.3.

As mentioned in the previous section, the ATLAS detector is a general purpose detector, with multiple physics goals:

- The search for and discovery of the Higgs boson. The Higgs boson has several decay channels, of which the best detectable channels are the decay to two

photons ($H \rightarrow \gamma\gamma$) for a light Higgs boson, while for a heavier Higgs boson the decay into two Z bosons or two W bosons is preferred, with their corresponding leptonic decay:

$H \rightarrow ZZ \rightarrow l^+l^-l^+l^-$ and $H \rightarrow W^+W^- \rightarrow l^+\nu l^-\bar{\nu}$. On the 4th of July 2012 it was announced that a Higgs-like scalar boson with a mass of ~ 125 GeV was found in both ATLAS [20] and CMS [21], which has since then been proven to be a Higgs boson [114].

- The search for direct evidence of physics beyond the standard model. Apart from searches for supersymmetry, this includes more exotic scenarios such as extra dimensions [115], micro black holes [116] or new heavy gauge bosons [117].
- Precision measurements of known SM particles, such as the top quark [118] and the W boson [119], and high precision tests of QCD [120].
- Studying the properties of a quark-gluon plasma from lead ion collisions [121].
- Flavour physics, studying CP-violation through the decay of B-hadrons [122].

Requirements

Each of these physics goals can be translated into requirements on the detectors. The detection of particles originating from Higgs decay for instance requires a good photon and lepton identification and energy (momentum) reconstruction, together with a good reconstruction of the missing transverse energy. A pair of produced supersymmetric particles would decay in cascades of particles, leading to high momentum jets, possible leptons and missing transverse energy, which would all need to be identified and reconstructed to a high precision. Thus a set of requirements has been formulated which the detector should meet:

- The high particle flux coming from the interaction point requires radiation hard detector elements with fine granularity, while the high collision rate requires fast electronics.
- A near hermetic coverage is needed to ensure no high momentum particles are missed. This is important for both the lepton and jet identification, and the reconstruction of both the transverse energy and missing transverse energy of an event.
- Good momentum resolution of charged particles is needed in the inner tracker system. The vertex reconstruction performed should be able to identify secondary vertices coming from τ or B-hadron decays.
- Good muon momentum resolution is needed for a wide range of muon momenta.
- Due to computing and storage limitations, a highly efficient trigger mechanism is needed on interesting events, with a high background rejection.

Detector component	Required resolution	η coverage
Tracking	$\sigma(p_T)/p_T = 0.05\% p_T \oplus 1\%$	$ \eta < 2.5$
EM calorimetry	$\sigma(E)/E = 10\%/\sqrt{E} \oplus 0.7\%$	$ \eta < 3.2$
Hadronic calorimetry (jets)	$\sigma(E)/E = 50\%/\sqrt{E} \oplus 3\%$	$ \eta < 3.2$
	$\sigma(E)/E = 100\%/\sqrt{E} \oplus 10\%$	$3.1 < \eta < 4.9$
Muon spectrometer	$\sigma(p_T)/p_T = 10\%$ at $p_T = 1$ TeV	$ \eta < 2.7$

Table 2.2: A summary of the performance requirements of the ATLAS detector. The resolution is parameterised as an energy (or momentum) dependent term and a constant term. For calorimeters this can be interpreted as follows: the first term is the stochastic fluctuations of the response, while the second constant term reflects local non-uniformities of the response, for instance due to leakage [107].

A quantitative summary of these requirements is shown in table 2.2, where requirements on the momentum and energy resolution of the various sub-detectors are given together with their spatial coverage, given as a range in the pseudorapidity η described in the following section. On top of this, a requirement on the resolution of the reconstruction of impact parameters and vertex positions exists. For high momentum particles in the central region of the detector, the design requires a transverse impact parameter resolution of $10 \mu\text{m}$ [123]. Here the impact parameter is defined as the closest perpendicular distance between a track and the primary vertex. In view of the lifetimes of $(c\tau)_\tau \sim 90 \mu\text{m}$ and $(c\tau)_b \sim 500 \mu\text{m}$ for τ leptons and B hadrons respectively [124], together with the corresponding Lorentz boost γ , the design requirement on the resolution of the main (primary) vertex is less than 1 mm in the z -direction [125].

In sections 2.3–2.7 each of the elements of the ATLAS detector is described in more detail, together with their operational performance. The reconstruction of objects is discussed in chapter 4, where also its performance is described, such as the momentum and energy resolutions.

Coordinate system

It is useful for the remainder of the thesis to specify the coordinate system used in the ATLAS detector. The nominal interaction point, right in the centre of the detector, is chosen as the origin. From here the x -axis points to the centre of the LHC ring, the y -axis points upwards and the z -axis lies along the beam-pipe. To obtain a right-handed coordinate system the z -axis points in anti-clockwise direction along the ring. The positive z -side of the detector is called side-A, whereas the opposite side is called side-C of the detector. The x - y plane is perpendicular to the beam direction and thus is referred to as the transverse plane. This plane is used often throughout this thesis, for instance for variables declared in this plane such as the transverse momentum

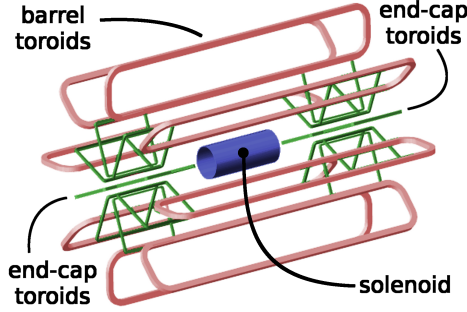


Figure 2.4: A schematic view of the ATLAS magnet system. Figure taken from [126].

$$p_T = \sqrt{p_x^2 + p_y^2}.$$

The cylindrical symmetry of the detector allows for the definition of several angles: the *azimuthal* angle $\phi \in [-\pi, \pi]$ is taken around the beam axis, and the *polar* angle $\theta \in [0, \pi]$ measures the angle with the beam axis. From this polar angle the more convenient *pseudorapidity* η is defined: $\eta = -\ln \tan(\theta/2)$ which is preferred above θ because the particle flux coming from the interaction is approximately constant as a function of η ¹. Furthermore, differences in η are invariant under Lorentz boosts. Finally the separation ΔR between objects in (ϕ, η) space can be defined as $\Delta R = \sqrt{\Delta\phi^2 + \Delta\eta^2}$.

2.3 Magnet system

The size and design of the ATLAS detector is mostly due to the choice of the magnet system.² Magnetic fields are imperative for the measurement of the momentum of charged particles, as these will deflect in the presence of such a field. The momentum of a particle determines the curvature of the track, hence measuring the curvature, when knowing the magnetic field, equals measuring its momentum. Particles travelling through the inner detector and the muon spectrometer, should therefore travel through a magnetic field in each sub-detector. These fields are produced by two magnet systems: a central superconducting solenoid magnet, providing a magnetic field inside the inner detector, and a superconducting toroid magnet system providing the field in the muon spectrometer. A schematic view of the magnet system is shown in figure 2.4.

¹To be precise, the particle flux is approximately constant as a function of the rapidity y , defined by $y = 1/2 \ln((E + p_z)/(E - p_z))$. For massless objects the pseudorapidity is very close to the rapidity. For massive objects, such as jets, this is not the case.

²As becomes obvious after close inspection of the name of the detector.

Solenoid magnet system

The cylindrical superconducting solenoid magnet is 5.8 m long, has a diameter of approximately 2.5 m and is embedded inside the electromagnetic calorimeter. The nominal current of 7730 A running through the bore provides a homogeneous 2 T magnetic field parallel to the beam axis inside the inner detector, which has a stored energy of 40 MJ. The design is optimised to keep the amount of material in front of the electromagnetic calorimeter as low as possible, adding up to only ~ 0.66 radiation lengths (X_0) on top of 1-2.5 X_0 of material in the inner detector. This is in part done by sharing its cryostat with the electromagnetic calorimeter.

Toroid magnet system

The toroid magnet system consists of a barrel toroid and two end-cap toroids, one on either side. All consist of eight superconducting rectangular coils placed in a cylindrical configuration – those of the barrel 25.3 m in length, with an inner diameter of 9.4 m and outer diameter of 20.1 m, while the end-cap coils are 5.0 m long and have an inner and outer radius of 1.65 m and 10.7 m, respectively. They thus define the size of the ATLAS detector. Operating at only 4.6 K, the average barrel field strength in the muon spectrometer is 0.5 T at a nominal current of 21 kA, storing an energy of 1.1 GJ. The end-cap toroids provide a local magnetic field in between the first two wheels of the muon spectrometer end-caps, with an average field strength of 1.0 T and a stored energy of 2×0.25 GJ.

2.4 The Inner Detector

2.4.1 General layout

The innermost subsystem, going by the appropriate name of ‘inner detector’, is a tracking system, designed to track the particles in their outward trajectory, measure their momentum and perform a vertex identification. To do this with high precision for the large number of particles traversing the detector at every collision, it consists of three complementary sub-detectors, from inside to outside: a silicon pixel detector, a silicon strip detector and a straw tube detector. These are arranged in three concentric cylinders around the beam axis, called the barrels, and on either side of each sub-detector a set of end-cap disks perpendicular to the beam axis. Figure 2.5 on the next page shows in the top figure a schematic view of one opened half of the ID and a schematic view of the barrel configuration in the middle figure. A more detailed blueprint of one quarter of the inner detector is shown in the bottom figure. As mentioned before, the requirements of the inner detector are to have a good momentum resolution for the pseudorapidity range of $|\eta| < 2.5$. The reconstruction of the high number of particles is accommodated by the fine granularity of the sub-detectors, while the momentum measurement is performed using the 2 T magnetic field provided by the solenoid magnet, which has been discussed in the previous section.

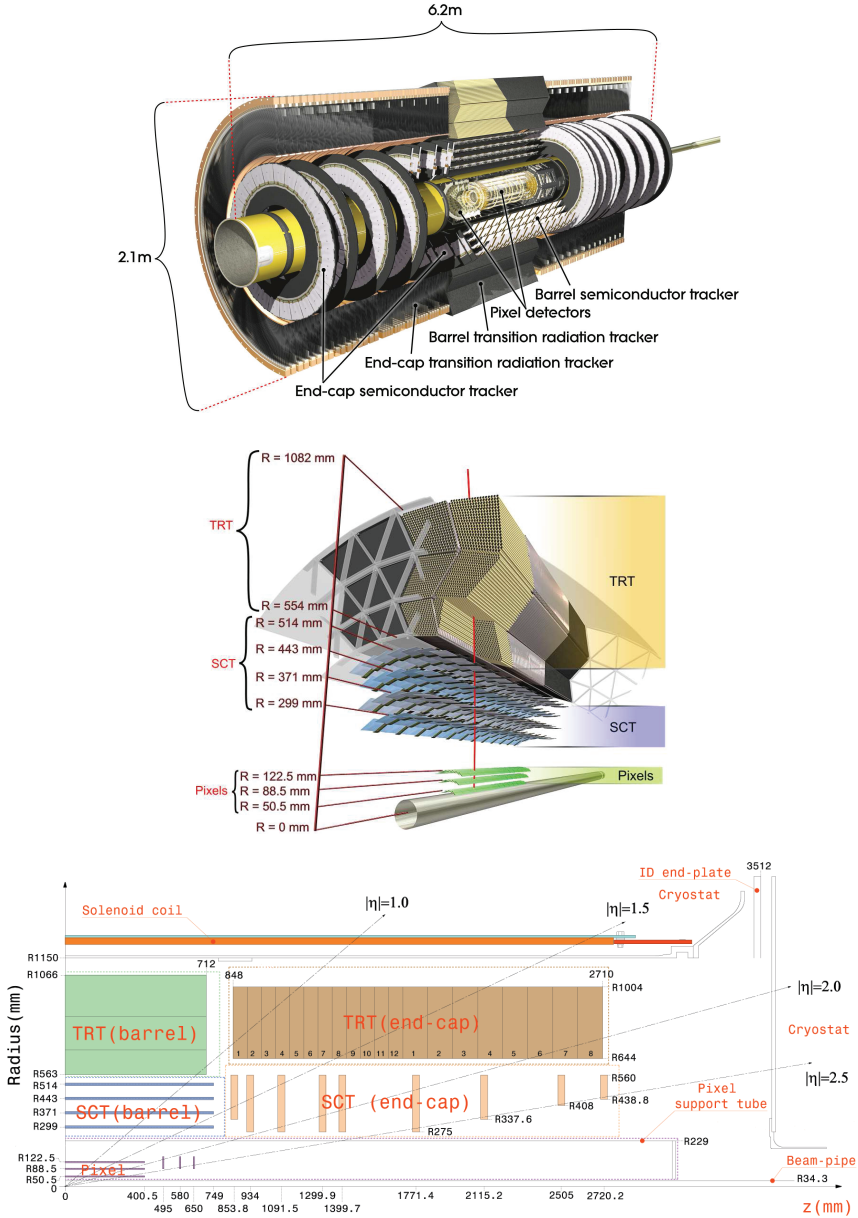


Figure 2.5: The ATLAS inner detector (ID). The top figure shows a cutaway view on the total inner detector, while the middle gives a better view on the dimensions and construction of the barrel. Taken from [107]. The bottom shows a schematic view of one quarter of the ATLAS inner detector, showing the configuration of the various barrel and end-cap layers of the three sub-detectors. Adapted from [107].

2.4.2 Pixel detector

The *pixel detector* is the innermost detector, and therefore subject to the highest particle flux. It should be therefore radiation hard, and have a fine enough granularity to separate the large number of tracks. It consists of three layers in the barrel with a radius of 50.5 mm, 88.5 mm and 122.5 mm, and three disks in each end-cap. This hermetic configuration assures that each particle typically crosses three pixel layers in its trajectory through the detector. The barrel and end-caps consist of a total of 1744 silicon modules, made of a 250 μm thick silicon wafer, implemented with readout pixels with a size of $50 \times 400 \mu\text{m}^2$, resulting in a total of 80.4 million readout channels. Particles traversing the pixel detector will ionise the charge depleted silicon layer³, resulting in a local charge deposition in a cluster of pixels. This charge deposition is read out by the dedicated electronics which is bump-bonded to the pixel sensor. The centre of the cluster is taken as the hit position, leading to a discrete space-point measurement. The spatial resolution of the detector is $10 \mu\text{m} \times 115 \mu\text{m}$ in $R-\phi \times z$ for the barrel, and the same in $R-\phi \times R$ in the end-caps. The difference between $R-\phi \times z$ and $R-\phi \times R$ comes from the configuration of barrel and end-caps respectively: in the barrel, modules are placed along the z -direction in a circular pattern around the beam pipe, while in the end-cap, modules are placed perpendicular to the beam pipe, along the R -direction. To protect the detector from radiation damage, the silicon sensors need to be cooled to approximately -5 to -10°C [127]. More details on the detection principle of silicon detectors is given in section 2.4.3, where the SCT is discussed in detail.

The barrel layer closest to the beam-pipe is essential for primary and secondary vertex reconstruction, and consequently the tagging of b -jets, and has the appropriate name B-layer. The proximity to the interaction point results in a high radiation environment. The original design required a replacement of the B-layer after three years of operation under nominal luminosity. This design has been revisited: in the two years of the first LHC shut-down (2013-2014) an extra pixel layer will be inserted, on the inside of the inner layer. This *insertable B-layer* [128,129] will improve the vertexing and tagging precision.

2.4.3 Semiconductor Tracker

The second high precision tracker is a silicon strip detector, called the *semiconductor tracker*, or SCT, adding four possible space-point measurements further from the interaction point to the track. Because of the larger distance from the interaction point, the SCT becomes increasingly important for the momentum determination when tracking particles with higher momenta.

Like the majority of ATLAS sub-detectors, the SCT is made up of a barrel and two end-caps – the former consisting of four concentric cylindrical layers, while of the latter each is made of 9 disks perpendicular to the beam axis at either side of the barrel. The barrel extends to $|\eta| < 1.4$, while the end-cap disks are placed such that

³The ionisation in the pixel detector is based on the same principle as that in the SCT, which is described in more detail in section 2.4.3

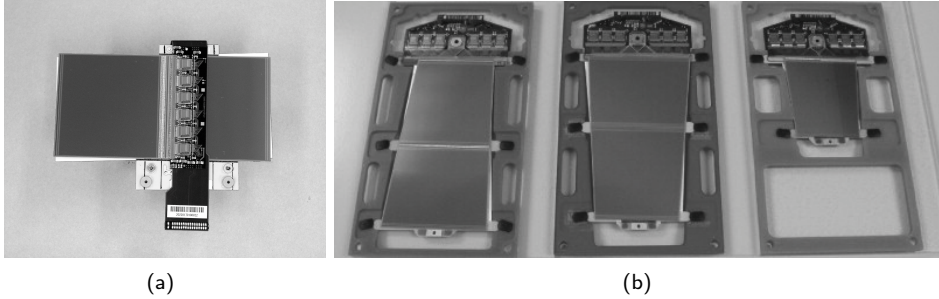


Figure 2.6: SCT modules, with a barrel module (a), and three end-cap modules (b), from left to right: an outer, middle and inner module. Note the angular offset between the upper and lower side of the module, and the hybrid with 6 chips. Adapted from [130] and [131].

the SCT provides four additional space-points for the total pseudorapidity range up to $|\eta| < 2.5$ as can be seen in figure 2.5 (c).

Each of the barrel and end-cap layers consists of smaller modules made up of two layers of $285\ \mu\text{m}$ thick silicon wafers, each containing 768 strips. In total the SCT consists of 4088 of such modules, with in total $63\ \text{m}^2$ of silicon, resulting in more than 6 million readout channels. In each barrel cylinder, one layer lies with the strips along the beam-axis, measuring the R - ϕ direction, as silicon strips do not give information in their longitudinal direction. The second layer has a small angular offset of $40\ \text{mrad}$ with respect to the first, thus enabling a measurement of the z direction for particles traversing both layers; likewise for the end-caps. The detection principle, explained in more detail below, is similar to that of the pixel detector. The choice for strips instead of pixels is due to the lower particle flux further away from the interaction point. The lower expected multiplicity requires less readout channels, reducing the cost of the detector. As is the case with the pixel detector, the silicon in the SCT needs to be cooled between approximately -5 to $-10\ ^\circ\text{C}$. Two module performance requirements are a detection efficiency higher than 99% and low noise occupancy, defined as the fraction of hits not related to a crossing particle, of less than 5×10^{-4} .

Each of the 2112 barrel modules (figure 2.6 (a)) has the same design: two $6.4\ \text{cm}$ long sensors are linked together by wire bonds, creating $12.8\ \text{cm}$ long wafers placed alongside one another. This results in a rectangular module with a width of $6.36\ \text{cm}$. Each side of the module has 768 readout strips with a high bias voltage applied, plus two inactive strips at either side defining the edge. These are tiled on top of each other, forming a hermetic coverage in the full ϕ range. The silicon strips in the barrel have a pitch (the distance between the middle of two neighbouring strips) of $80\ \mu\text{m}$.

The end-cap design is somewhat different, with disks made of three rings of trapezium shaped modules. The inner, middle and outer rings are made of different types of modules, with a different geometrical design to achieve the best hermetic acceptance. Due to these geometrical reasons, disk 9 (the furthest from the interaction point) has

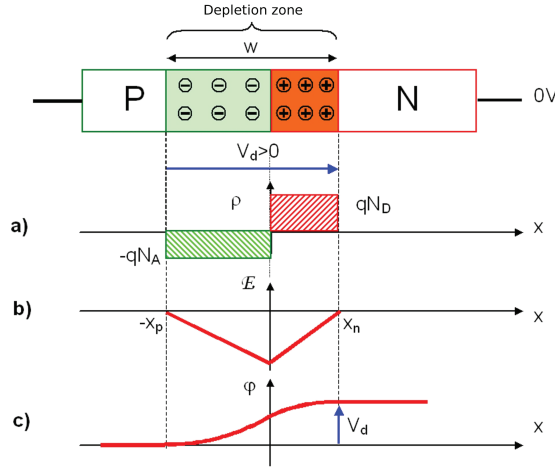


Figure 2.7: Diagram of a *p-n* junction at zero bias voltage. The top shows the *p*-type silicon on the left and *n*-type on the right. The three diagrams below show the charge density (a), the electric field (b) and the potential (c) in the depletion region. Taken from [132].

only outer modules, while disks 1, 7 and 8 have no inner modules. Furthermore, the inner modules have only one set of sensors, making them half as long – the same is true for the middle modules of disk 8, which are thus called ‘short middle’ modules. This configuration can be seen in figure 2.5. Due to the trapezium shape, the pitch of the strips vary in the end-cap, from $56.9 \mu\text{m}$ to $94.2 \mu\text{m}$.

The two layers of the module, with the small angular offset, are glued on top of a thin base-board made of thermal pyrolytic graphite, providing the structure and thermal conductivity path between sensors and coolant. Each of the sensors is read out by a Application Specific Integrated Circuits (ASIC) chip, with 6 chips per module side. Each chip thus operates 128 readout channels. These chips are mounted on a hybrid, placed in the middle of the barrel modules and at the end of the end-cap modules, as seen in figure 2.6.

Detection principle of the SCT

Semiconductor devices build upon the principle of *doping*, contaminating the silicon with impurities such that the conductivity is changed. Doping silicon can lead to *p*-type material, with an excess of holes, or *n*-type material with an excess of electrons. A ‘hole’ in this sense is nothing else than the absence of an electron in the lattice of silicon atoms, which can travel through the lattice as if it were a positively charged particle. Bringing *n*-type and *p*-type doped silicon in contact with each other, electrons move from the *n*- to the *p*-type, and inversely for the holes, creating a so called *p-n* junction. Near the contact layer of the two materials, a *depletion zone* is created, with an absence of free charge carriers. Yet the electrons diffusing to the *p*-type material

leave positively charged ions in the n-type region, and vice versa for the holes. Thus two oppositely charged layers are created near the p-n interface, creating an electric field. As this electric field is in the opposite direction to the diffusion current, it stops the flow of the current, arriving at a static situation. This is illustrated in figure 2.7. When an external *reverse bias* potential is applied to the silicon, with the p-type silicon attached to the negative terminal, the already built up potential is increased, increasing the depletion zone. In a detecting device, this depletion zone is ideally the full thickness of the sensor.

When a charged particle travels through semiconducting material, electron-hole pairs are produced – a minimum ionising particle travelling through a $285\text{ }\mu\text{m}$ thick layer of silicon will create approximately 25000 electron-hole pairs, generating a charge deposition of $\sim 4\text{ fC}$ in the silicon [133]. If this happens in a depletion zone, the holes will drift to the p-side, while electrons drift to the n-side of the material. Here the charge is collected on metal readout strips, with a signal strength proportional to the induced current. The SCT uses for the silicon wafers an n-type bulk, with p-type strips implanted. On top of these implants, resting on an insulator, metal readout strips are placed with a nominal bias voltage of 150 V [127].

The strip pitch defines the tracking resolution: for a finer resolution, a lower strip pitch should be achieved. The intrinsic resolution of SCT modules is $17\text{ }\mu\text{m}$ in the $R\text{--}\phi$ direction and $580\text{ }\mu\text{m}$ in the z or R direction, for barrel or end-cap modules respectively.

SCT readout

Approximately 10 ns after a particle traverses the silicon, the charge is collected by the readout strips [127]. The current is amplified by two amplifiers on the corresponding chip. A discriminator then compares the amplitude of the signal against a pre-defined threshold. This threshold is taken nominally at 1 fC , which allows the detector to reach its performance goals with respect to the noise occupancy and efficiency [127]. The output of the discriminator is either a '1', when the signal is above threshold, or a '0' when it is not. For each of the 128 channels of a chip this results in a binary 'hit or no hit' decision, stored as a single bit in a 132-bit pipeline. This simple binary decision allows for a simple chip, reducing the size and power consumption and thus the cost.

When the readout buffer receives the level-1 trigger of a certain bunch crossing (see section 2.7), the corresponding data is passed to the data compression logic, together with that of the previous and following bunch crossing. The compression logic only transfers the data of an event on to the readout logic if the channel obeys the required hit pattern. Four of these hit patterns are available, shown in table 2.3. The *Any hit* pattern requires at least a hit in one of the three available bunch crossings, without requirement on hits in the other bunch crossings. This pattern is used for detector studies. *Level sensing* and *edge sensing* are the two patterns used in data taking. The former requires at least a hit in the bunch crossing corresponding to the L1 trigger, and is used for regular data taking. The latter, edge sensing, requires no hit in the previous bunch crossing, a hit in the correct bunch crossing and anything in the

Selection criteria	Hit pattern	Usage
Test	XXX	Test mode
Any hit	1XX, X1X or XX1	Detector alignment
Level sensing	X1X	Data taking
Edge sensing	01X	Data taking

Table 2.3: *The available hit patterns for the compression logic. The middle column represents the three time-bins read out: from left to right the previous, the current and the following bunch crossing. A 1 represents a hit in the corresponding time-bin, a 0 no hit, and an X means no requirement. Adapted from [127].*

following crossing. It is designed for data taking with 25 ns LHC bunch spacing. The pattern can be explained as follows: one would not expect a hit in the strip before the bunch crossing itself, yet as the length of the signal might be longer than 25 ns, leading to a signal over the threshold for two bunch crossings, there might a hit in the following time-bin. The readout logic finally reads the data, and transmits it from the chip via two optical links to off-detector electronics. The chips are chained, creating redundancy such that malfunctioning chips can be bypassed.

2.4.4 Transition Radiation Tracker

The third ID sub-detector, furthest from the beam-pipe, is a straw tube detector called the *transition radiation tracker* (TRT). Its design provides the tracking algorithm with many space-points, while also identifying electrons through transition radiation. The tracking is done with 73 layers of straw tubes in the barrel, and 160 layers in the end-caps.

The straw tubes, with a diameter of 4 mm and a length up to 144 cm, are filled with a xenon-based gas mixture, with a gold plated tungsten anode wire running through the middle of the tube. The potential difference between tube and wire cause free electrons created by passing charged particles ionising the gas to drift towards the wire. The drift time is proportional to the distance of the traversing particle to the wire.

These straws are embedded in polypropylene fibres and foil in the barrel and end-caps respectively, with varying refractive indices. Relativistic charged particles traversing transitions in the dielectric constant will emit X-ray photons with an intensity that is directly proportional to their Lorentz factor γ . The photons are absorbed by the xenon in the straw tubes, where, as the electron mass is 273 times smaller than the charged pion mass, electrons will produce a higher signal than pions for the same momentum. Applying two thresholds, one optimised to detect ionisation and one optimised to detect transition radiation from electrons, the TRT provides a method to distinguish electrons from pions.

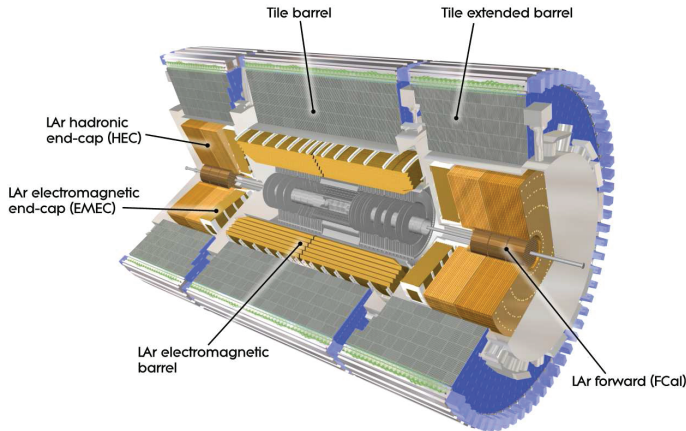


Figure 2.8: *Cutaway of the ATLAS calorimeter. The electromagnetic calorimeter is shown in yellow, while the hadronic tile calorimeter is shown in grey. The cracks between the barrel and extended barrel or end-caps are clearly visible. Taken from [107].*

2.5 Calorimeters

The second set of sub-detectors a particle produced at a collision encounters are the calorimeters. These measure the energy of both charged and neutral particles, and are the only detector element able to identify the latter. The ATLAS calorimeter system consists of two types of calorimeters: closest to the interaction point is the electromagnetic calorimeter, designed to measure the energy of electrons and photons through their electromagnetic interactions; further from the interaction point lies the hadronic calorimeter, designed to detect showers from hadrons and hadronically decaying taus, and measure their energy through both their electromagnetic and strong interactions. In the barrel the calorimeter system is cylindrical and extends radially from approximately 1.4 m to 4.2 m from the interaction point. It consists of a barrel and two end-cap regions, which together cover the pseudorapidity range up to $|\eta| < 3.2$, plus an additional forward detector on either side providing coverage for the range $3.1 < |\eta| < 4.9$. A view on the calorimeter system is shown in figure 2.8, where a part is cut out for a better view of the inside.

ATLAS uses *sampling* calorimeters, which consist of multiple layers of alternating an absorbing material and an active medium. A particle interacting with the dense absorbing material loses some of its energy in the production of particle showers. Electromagnetic (EM) showers produced by electrons and photons are the result of Bremsstrahlung and e^+e^- pair production from the interaction with the absorbing material. Hadronic showers on the other hand are caused by interactions with nuclei of the hadron with the absorber, resulting in a shower of low energy hadrons (e.g.

pions). The subsequent layer containing the active medium measures the energy deposited by charged particles in these showers. After several of these sequences the incoming particles will have lost all their energy. The number of produced shower particles will be proportional to the energy of the initial particle, thus by measuring the total shower the energy of the incoming particle can be calculated (after careful calibration).

Charged particles moving through a medium lose energy by ionisation and photon emission (bremsstrahlung). The energy-loss via ionisation depends via the Bethe formula [19] on the properties of the active medium, but also on the particles speed and mass – a lighter particle loses more energy. The higher mass of a muon with respect to an electron requires a much larger stopping length for the muon to lose all its energy, which means muons with a momentum above ~ 3 GeV will not be stopped in the calorimeter and will reach the muon spectrometer (see section 2.6). On the other hand, ultrarelativistic charged particles lose more energy via bremsstrahlung with increasing $\beta = v/c$. The high velocity of high-momentum electrons ensures that these lose all of their energy in the calorimeter. The weakly interacting neutrino has an energy dependent cross section for nuclear interactions of only $\sim 10^{-38}$ cm² [134]. It typically does not interact with the calorimeter (or any other detector system) and thus is not measured. However, some information can be found from the *missing transverse energy* or $\cancel{E}_T$, which is described in detail in section 4.5.

Electromagnetic Calorimeter

The electromagnetic calorimeter (ECAL) consists of a barrel section, constructed as a cylinder around the interaction point extending to $|\eta| < 1.475$, and an end-cap wheel at either side spanning the pseudorapidity range $1.375 < |\eta| < 3.2$. These both use the same design, using lead plates as absorbing material and liquid argon (LAr) as the active material. The lead plates have an accordion shape, giving a gap-less ϕ coverage. The choice for LAr in between these plates is supported by its linear response and its radiation hardness. The signal, in the form of a current from free charges in the LAr, is read out by electrodes between the absorbers.

The design of the ECAL is optimised to stop all electromagnetic (EM) showers, with maximum possible transparency for hadronic objects such as jets. This is achieved with an ECAL thickness corresponding to between 22 and 33 radiation lengths (X_0) for the barrel and between 24 and 38 in the end-cap, while the material amounts to only ~ 1.5 nuclear interaction lengths λ . Here X_0 is the characteristic interaction length in the material, defined by the distance it takes an electron to lose all but a factor $1/e$ of its energy by Bremsstrahlung, while λ is the mean distance travelled by a hadron before interacting with a nucleus.

The barrel calorimeter is constructed in three layers: the innermost layer, which is the smallest with a depth of $4.3 X_0$, has the finest segmentation in η with a granularity of $\Delta\eta \times \Delta\phi = 0.003 \times 0.1$, allows for a better separation of photons and π^0 versus electrons. The second layer has a coarser η granularity at $\Delta\eta \times \Delta\phi = 0.025 \times 0.025$ and makes up for the largest part of the calorimeter with a depth of $16 X_0$. The last layer,

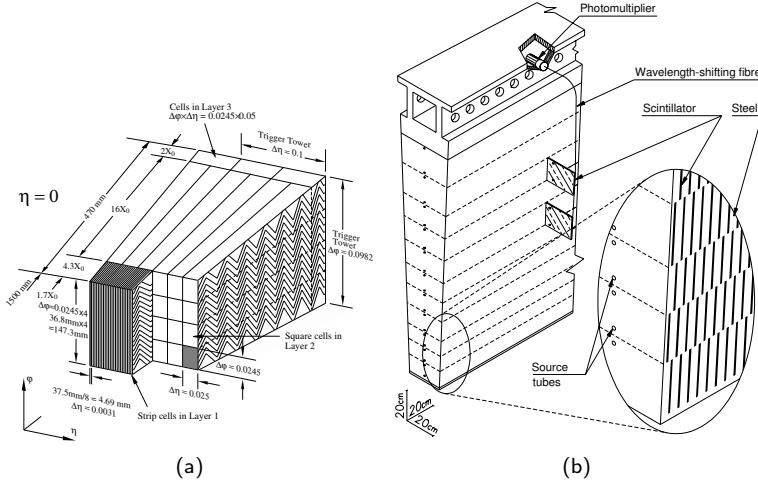


Figure 2.9: Schematic view of an electromagnetic calorimeter module (a) and tile calorimeter module (b). The three layers with varying granularity in the ECAL module are indicated, just as the accordion shaped lead absorber. The schematic view of the HCAL module shows the varying layers of steel absorber and scintillator, as well as the photomultiplier configuration. Taken from [107].

of only $2 X_0$ in depth, has again a coarser granularity. Its purpose is to distinguish EM showers from hadronic showers, which have most of their energy deposit at a larger distance. This configuration can be seen in figure 2.9 (a). For the largest part, the end-cap wheels have the same specifics, with a slightly deeper detector. The innermost part, with $2.5 < |\eta| < 3.2$, only has two coarser grained layers.

The space between the electrodes and absorbers is 2.1 mm , corresponding to a drift time of 450 ns . This is significantly longer than the bunch spacing of the LHC, and thus incorporates information of previous collisions, leading to out-of-time pile-up. This is countered by using a bipolar signal shaper, which quickly produces a short signal pulse proportional to the peak current [135].

The ‘crack’-region between the barrel and end-caps at $1.37 < |\eta| < 1.52$ has a degraded energy resolution, and is thus not used for precision measurements.

Hadronic Calorimeter

Just like the ECAL it surrounds, the hadronic calorimeter (HCAL) is constructed as a cylindrical barrel and two end-cap wheels. Its barrel system consists of a central barrel spanning $|\eta| < 1.0$, and an extension on either side covering the region $0.8 < |\eta| < 1.7$. The gap between these is used for services for the inner detector, EM calorimeter and

solenoid magnet.⁴ Both barrel and end-caps have an inner radius of 2.3 m and an outer radius of 4.3 m. This sampling calorimeter uses steel as absorber, and scintillating plastic tiles as active material – the hadronic barrel calorimeter therefore is also known as the *tile calorimeter*. Particle showers entering the tile calorimeter produce light in the scintillating tiles, which is collected and transported to photomultiplier tubes which amplify and read out the signal. As hadronic showers are wider than EM showers, the granularity of the tile calorimeter can be coarser than that of the ECAL, at 0.1×0.1 in $\Delta\eta \times \Delta\phi$. A schematic view of a HCAL module can be seen in figure 2.9 (b). The detector response time of the tile calorimeter is approximately 75 ns.

The hadronic end-cap calorimeters, HECs, are again liquid argon calorimeters, consisting of two wheels on each side which are both made of parallel copper plates with LAr as active material between them. The wheels cover a range of $1.5 < |\eta| < 3.2$, and are built of 32 wedge shaped modules. The inner wheel's copper plates are 25 mm thick, while those of the outer wheel are twice that. Its detection principle is the same as for the EM calorimeter discussed before.

The depth of the hadronic calorimeter is approximately 7.5λ for the barrel, with λ the nuclear interaction length, and 10λ for the end-cap. This is enough to contain most energetic showers, with low probability of a *punch-through* to the muon system. In such an event the hadronic shower is not stopped completely by the calorimeter and spills over into the muon spectrometer.

Forward Calorimeter

For the determination of the missing transverse energy, it is important to have a good hermetic coverage. Therefore, a separate calorimeter is constructed which can withstand the high particle fluxes in the very forward region, $3, 1 < |\eta| < 4.9$, which comes down to a radius of only ~ 8 cm around the beam. This forward calorimeter (FCAL) consists of three wheels on either side, which are all liquid argon-based. The innermost is optimised for electromagnetic showers and has copper absorbers, while the outer two measure the hadronic showers and have tungsten as the absorbing material. The high particle flux requires a large depth of the calorimeter: the wheels have a depth of $27.6 X_0$, $91.3 X_0$ and $89.2 X_0$ from inner to outer wheel respectively.

2.6 The Muon Spectrometer

The muon spectrometer is the outermost system of ATLAS. Of all possible SM particles produced by the collision or in subsequent decays (except neutrinos), the muons will travel the furthest, due to their long lifetime and low energy emission in the form of Bremsstrahlung, and are not easily stopped by the calorimeters. The muon system identifies muons, and measures their position and momentum. The muon momentum can be measured by the muon spectrometer alone for muons with

⁴The given pseudorapidity ranges causes one to think there is not a gap, but an overlap between the central barrel and the extension. Note that this is caused by the definition of the pseudorapidity, which is an *angle*. Figure 2.8 shows the gap between the central barrel and its extension.

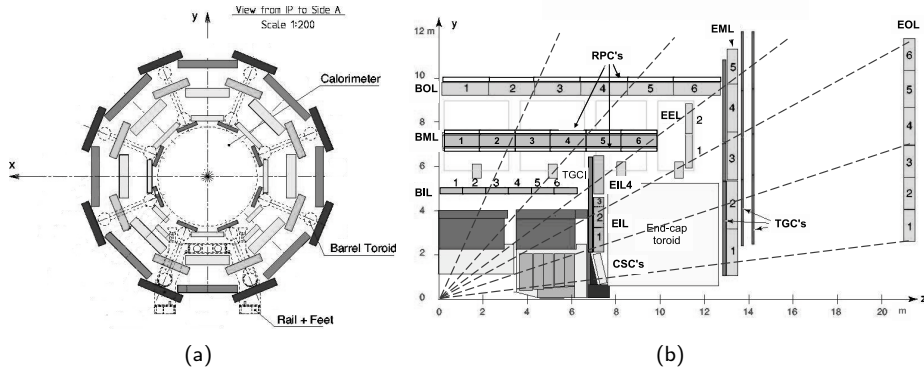


Figure 2.10: Cross section of the muon spectrometer: (a) in the x - y plane, with the interaction point in the middle, and (b) of one quarter of the y - z plane. In the left figure the overlap in the placing of the various MDT stations is clearly visible. In (b), the stations denoted with a number are all MDTs. The dashed lines denote the trajectories muons with infinite momentum would follow, crossing three muon stations each. Taken from [107].

a momentum between about 3 GeV and 3 TeV, although the momentum resolution for low momentum muons is driven by the inner detector. It is designed to have a momentum resolution of 10% for muons with a p_T of 1 TeV.

Precision tracking

The momentum measurement is performed by precisely measuring the track of a particle in a strong magnetic field. From the curvature of the track, and knowing the magnetic field strength, the momentum can be deduced. High momentum tracks will have smaller curvature, therefore both a strong magnetic field and a large distance between the position measurements is needed for an accurate momentum measurement. This drives the construction of the muon spectrometer, shown in figure 2.10.

The precision tracking is performed by two detection systems. *Monitored drift tubes* (MDTs) are placed in three layers in the barrel region, $|\eta| < 2.0$, at approximately 5 m, 7.5 m and 10 m from the interaction point. The layers consist of 16 MDT chambers overlapping slightly, as seen in figure 2.10 (a). There is a hole in this configuration at $\eta = 0$, to provide for services to the underlying detectors and solenoid magnet. This gap, which spans an angular range of $|\eta| \leq 0.08$ at its largest, leads to a resolution degradation, discussed in section 4.3.1. The forward region, $2.0 < |\eta| < 2.7$, is built up of three wheels, of which the outermost two are made up of solely MDTs, while closest to the beam-pipe the inner wheel has a layer of *cathode-strip chambers* (CSCs), which are better adapted to the high particle flux in this region. An additional gap in the coverage is caused by the ‘feet’ of the detector support structure, in sectors 12 and 14 (see figure 2.10 (a)).

The MDTs are configurations of 30 mm wide aluminium tubes, between 1 and 6 m

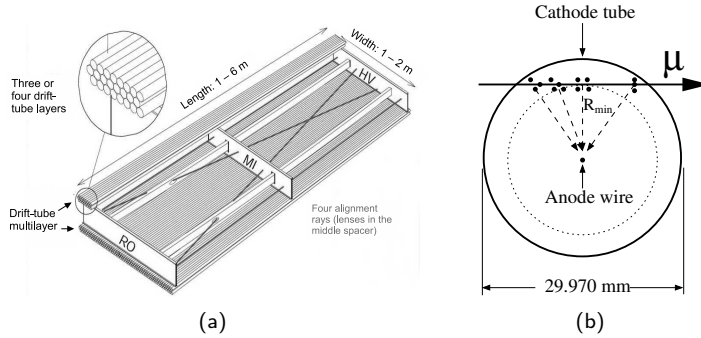


Figure 2.11: Schematic view of an MDT station, with three of four drift tube layers stacked on top of each other (a) and cross section of a drift tube (b), with the detection principle shown.

in length, filled with Ar/CO_2 gas at 3 bar, with a tungsten-rhenium anode wire in the centre on which a potential of 3080 V is applied. A traversing charged particle will ionise the gas – the freed electrons will free new electrons, creating an avalanche of electrons – see figure 2.11. This causes an electrical signal in the wire, which can be read out. Each MDT chamber consists of three to eight layers of drift tubes. As the tubes are placed along the ϕ direction, the MDTs only provide a high precision in the R - z plane. Unlike the MDTs, the CSCs have two cathode strips per plane, placed perpendicular to each other, and multiple anode wires, making them multi-wire proportional chambers. There are four of these planes in each CSC chamber, which are trapezium shaped, measuring ~ 1.1 m in length and $\sim 0.5 - 1$ m in width [136]. The charges induced on the cathode strips are read out instead of the wire.

The bending of the muons is achieved by the large toroid magnet for pseudorapidities below 1.4, while it is achieved for $|\eta| > 1.6$ by two smaller end-cap magnets on either side. For the transition region between these, $1.4 < |\eta| < 1.6$, a combination of the barrel and end-cap fields is used. This configuration gives approximately an orthogonal field to trajectories of muons from the interaction point, as discussed in section 2.3.

Trigger chambers

The triggering on muons is done by *resistive plate chambers* (RPCs) in the barrel region ($|\eta| < 1.05$) and *thin gap chambers* (TGCs) in the end-cap wheels, where $1.05 < |\eta| < 2.4$. To be able to trigger on muons, these systems should provide fast information on the tracks inside the detector, with information on both azimuthal and radial components. Aside from the trigger function of both trigger systems, the information on the azimuthal coordinate complements the MDT measurement, which can only be done in the R - z direction. The aforementioned gaps in the MDT coverage due to the services and support structure occur in the trigger chambers as well, leading to a degradation of the triggering efficiency at these coordinates.

The conditions differ between the barrel and end-cap regions, with higher radiation

levels in the end-cap region, while the $|\eta|$ dependence of the muon p_T requires a finer granularity in the end-cap to match the p_T resolution in the barrel. These requirements lead to trigger chambers with different designs between the barrel and end-cap.

Two RPC stations are placed below and above the middle MDT layer, while a third is placed on the outer side of the outer MDT layer. Each station consists of two RPC layers, corresponding to a total of six measuring points of η and ϕ . This redundancy improves the efficiency and helps to reject fake tracks, coming for instance from detector noise. The large distance between the outer and inner chambers allows for triggering on high momentum tracks, with a lower thresholds from 9 to 35 GeV, while the small distance between the two innermost chambers extends this to low p_T (6 – 9 GeV) [137]. The RPCs are gaseous detectors, which unlike the previously described chambers are wireless. An RPC consists of two parallel resistive plates placed 2 mm apart, with an electric field of 9.8 kV between them. Charged tracks traversing the detector ionise the gas, and the liberated electrons will accelerate through the electric field, causing avalanches of charged particles to move towards the anode. This is detected by readout strips mounted on the anode. The RPCs are designed to have a spatial resolution of 5-10 mm, and a time resolution of ~ 1 ns [137].

In the end-cap, the triggering is performed by TGC stations, one of which is placed on the inner layer of the end-cap wheels, while three surround the middle MDT wheel. They are based on the same multi-wire principle applied in the CSCs, yet unlike these the TGCs have a smaller distance between the wires (1.4 mm) than between wire and anode (1.8 mm). This guarantees drift times shorter than 25 ns, the time between two consecutive bunch crossings. The radial readout strips provide a measurement on the ϕ coordinate, while the η coordinate is measured using the anode wires. Again each station consists of two TGC layers except the outermost station, which only has one. The lack of TGCs at the outer MDT wheel is no problem: as there is no magnetic field between the outermost wheels, the extrapolation towards the outer layer can be done accurately.

2.7 Trigger system

After every collision, the particles coming from the interaction point are detected by one or more ATLAS sub-detectors. The electrical signals coming from their respective readout modules are subsequently stored. As discussed at the beginning of the chapter, there is a collision every 25 ns under nominal LHC operations, leading to a collision rate of 40 MHz. However, not all these events can be stored: at an the average event size of 1.6 MB, the high collision rate would result in storing 60 TB every second, which is unfeasible in terms of both computing power for event reconstruction and data storage requirements. At the same time, the majority of the collisions are uninteresting ‘soft collisions’, where the centre-of-mass energy of the interacting partons is not enough to produce particles corresponding to the physics goals of ATLAS.

To solve this problem, ATLAS possesses a highly selective trigger system which pre-selects interesting events as they happen, bringing down the event rate by a factor of 10^5 to a value of ~ 400 Hz. This final rate is mostly limited by the available storage

space.

2.7.1 Three trigger levels

The trigger system consists of three consecutive levels, which select events based on certain criteria: level 1 (L1), level 2 (L2) and event filter (EF). Each following level has a refined set of criteria, bringing the trigger accept rate down.

Level 1 trigger

The L1 trigger performs a hardware-based selection, using reduced-granularity information from the calorimeters and muon trigger chambers (RPC and TGC), in less than $2.5 \mu\text{s}$. The L1 calorimeter trigger (L1Calo) identifies large energy deposits in the calorimeters, coming from electrons, jets or photons, as well as events with large E_T . Furthermore, isolation requirements can be imposed on electrons and photons, setting a limit on the total deposited energy in a ring around the region of interest (RoI), giving the coordinates in (η, ϕ) where interesting features are seen. Likewise a threshold is placed on the hadronic energy behind an EM cluster. Several energy thresholds can be programmed for each object type, above which the L1Calo fires. The muon L1 trigger searches for signals from both the RPC and TGC, corresponding to a track coming from the interaction point. Events with hit patterns consistent with a high p_T muon are used, if the muon is above one of six programmable muon p_T thresholds.

The final decision is made by combining the information of both the L1Calo and L1 muon trigger, and comparing to 256 items in the trigger menu. If an event contains any signature in the menu, the event is passed on to L2, together with RoIs. The L1 trigger reduces the event rate from 40 MHz to below ~ 75 kHz. With increasing instantaneous luminosity the triggers using low p_T thresholds need to be prescaled to keep the rate at the required value. This also holds for the next two levels of the trigger.

Level 2 trigger

The regions of interest defined by the L1 trigger are further scrutinised by the L2 trigger. This software-based trigger uses information from the full detector only within the RoIs to do a fast reconstruction with finer detector granularity, while refining the selection criteria. Because of the reconstruction it takes longer than the L1 trigger, at 40 ms per event. This trigger reduces the event rate to ~ 6 kHz, which depends on the available computing power.

Event filter

The last in the series is the event filter (EF), which looks into the full reconstructed events passing L2. This stage takes approximately 4 s per event, requiring a computing farm with approximately 1500 computers to reduce the rate to 400 Hz. Events passing the EF are moved to permanent event storage.

2.7.2 Trigger menu and streams

The total sequence of algorithms an event passes is called the *trigger chain*, and is used to select events in offline analyses. The naming of these trigger chains tells one about the triggers the event has passed – for instance, an event passing EF_mu24i_tight has one isolated muon which breaches the momentum threshold of 24 GeV, and passes a ‘tight’ muon criterion. ATLAS runs with many of these trigger chains simultaneously in a trigger menu (up to 1000 in 2012). These need to be updated with increasing instantaneous luminosity to be able to still fulfil the requirement on the event rate.

The trigger system selects events for various purposes, the most apparent being physics analysis, but also other purposes such as calibration and monitoring. The EF not only performs the last event selection, but also categorises the event. For further physics analysis the events are categorised in three physics streams, according to their significant objects: the Muons, Egamma and JetTauEtmiss stream. For instance, an event containing an electron and two high p_T jets will be added to both the Egamma and JetTauEtmiss stream. In this thesis, focussing on the search for supersymmetric in fully hadronic events, we will use the JetTauEtmiss stream most often. A fourth physics stream, MinBias, contains a random sample of events. These minimum bias events allow for a search in events where no interesting features were expected. Furthermore, the stream can be treated as a background sample, for instance for detector performance studies.

2.8 Operational performance of the ATLAS detector

The performance of the ATLAS detector has been outstanding. During the three years of operation, ATLAS has recorded data with an average efficiency of 93.2%, while each of the sub-detectors has recording data $> 99\%$ of the time in stable beams. The operational status of the ATLAS detector in the middle of October 2012 is shown in table 2.4. It should be noted that this is a snapshot, and the precise operational fraction can change from run to run. This section discusses the operational performance of the various sub-detectors. As mentioned before, the performance of the reconstructed objects can be found in chapter 4. First the hit efficiency of the SCT sub-detector will be discussed, after which the operational status of the other sub-detectors is addressed.

2.8.1 Operational status and hit efficiency of the SCT

To ensure high quality tracking, a high efficiency of all tracking sub-detectors is imperative. In the SCT, this efficiency consists of three components: the operational status, describing the relative number of operational SCT silicon components during a run; the data-taking efficiency, representing the relative time the detector elements are turned on and can send back information; and an intrinsic hit efficiency, giving how well the silicon sensors which are operational are performing.

Sub-detector	Number of channels	Approximate operational fraction
Pixels	80×10^6	95.0%
SCT Silicon Strips	6.3×10^6	99.3%
TRT Transition Radiation Tracker	3.5×10^5	97.5%
LAr EM Calorimeter	1.7×10^5	99.9%
Tile calorimeter	9.8×10^3	98.3%
Hadronic end-cap LAr calorimeter	5600	99.6%
Forward LAr calorimeter	3500	99.8%
LVL1 Calo trigger	7160	100%
LVL1 Muon RPC trigger	3.7×10^5	100%
LVL1 Muon TGC trigger	3.2×10^5	100%
MDT Muon Drift Tubes	3.5×10^5	99.7%
CSC Cathode Strip Chambers	3.1×10^4	96.0%
RPC Barrel Muon Chambers	3.7×10^5	97.1%
TGC End-cap Muon Chambers	3.2×10^5	98.2%

Table 2.4: *The operational status of the ATLAS detector on the 19th of October 2012.*

SCT component	Barrel	End-cap A	End-cap C	Total SCT	Fraction
Modules	10	4	14	28	0.68%
Chips	35	6	5	46	0.09%
Strips	3859	3501	3787	11147	0.18%

Table 2.5: *Number and fraction of disabled modules, chips and strips in the SCT on May 20th 2012. Note that the total number of modules in the barrel and each of the end-caps is 2112 and 998, respectively, with each 12 chips and 1536 strips. The number of chips (strips) exclude those which are automatically disabled on disabled modules (modules and chips).*

Operational status and data-taking efficiency

During the first run of the LHC (Run 1), spanning three years from 2010 to the end of 2012, the SCT has been operating very stably. In every data taking run at least 99% of all SCT modules were operational, while in general less than 0.1% individual chips were disabled, and less than 0.2% individual silicon strips. Table 2.5 lists the number of disabled detector components in May 2012. Note that, although these numbers are typical for Run 1, they represent just a snapshot, with the precise numbers varying slightly over time. Half of the disabled modules are due to a leak in the cooling device

for one quadrant of the outermost disk of end-cap C. All modules cooled by this cooling loop had to be disabled. The remaining half of the disabled modules are predominantly due to malfunctioning connections on the detector side. Disabled strips are mainly due to constant high noise. Finally, disabled chips are mostly due to failing optical read out links: when an optical link fails, the link on a neighbouring module can take over. This rerouting method does however lead to the information loss of one chip. The number of optical links using this redundancy was around 1% at the end of Run 1.

The SCT has been fully operational during the data taking periods in Run 1. It recorded 99.9%, 99.6% and 99.1% of the proton-proton data delivered by the LHC in 2010, 2011 and 2012, respectively. The small losses are predominantly due to errors in the read-out driver boards (RODs), which perform the control and data handling of up to 48 modules each. When a ROD gave such an error, a signal ('busy') was sent to the data acquisition system inhibiting data taking by the modules of that ROD. Apart from this error, a potential source of data loss comes from the time it takes between the declaration of stable beam conditions by the LHC operators and increasing the bias voltage from 50 V at standby to the nominal bias voltage for of 150 V on all SCT modules. However, as the turning on of the SCT typically takes 60 seconds, less than several other sub-detectors, the data lost due to this effect is negligible.

SCT hit efficiency definition

With the highest luminosity ever achieved, the effect of radiation damage and the so-called *ageing* of the detector is very important to keep in check. One of the main properties of the SCT which is influenced by this effect is the intrinsic hit efficiency.

The hit efficiency of the SCT is computed using a 'holes on track' method, where for each well-reconstructed track the extrapolated trajectory through the detector is scrutinised. For every possible encounter of the track with an SCT module the number of clusters or hits are counted, as well as the number of 'holes'. A hole is defined as a lack of a hit where one is expected. Here the inactive material around each module is taken into account: if the intersection of the track with the SCT is outside the sensitive area of the module or within 3 standard deviations from the edge of the active material, it is not counted as a hole. This procedure ensures we measure the efficiency of the active material. Likewise, inactive modules and chips are excluded from the calculation; however, isolated inactive strips are not excluded, and thus contribute to the inefficiency. The hit efficiency ϵ_{hit} thus tells us how often a hit is detected when there should have been one, and is defined as:

$$\epsilon_{\text{hit}} = \frac{N_{\text{hit}}}{N_{\text{hit}} + N_{\text{hole}}}, \quad (2.1)$$

with N_{hit} the number of found hits and N_{hole} the number of holes on the track.

Misalignments of the SCT might lead to a bias in the track fitting procedure, affecting the calculated efficiency. To reduce this bias, hits are added to a track if they are within 200 μm of a track intersection. Only 4.4% of these hits are due

Track collection	COMBINEDINDET Tracks	SCTSTANDALONE Tracks
Number of other hits	≥ 6	≥ 7
χ^2 cut	< 3	< 3
$ d_0 $ cut	< 10 mm	< 10 mm
p_T cut	> 1 GeV	> 1 GeV
Number of tracks	< 250	< 250

Table 2.6: *Track selection for COMBINEDINDET tracks and SCTSTANDALONE tracks used in the SCT hit efficiency calculation.*

to noise or incorrect track associations, while the remainder indeed originates from track reconstruction inefficiencies [138], where the tracking algorithm has wrongly not included the hit in the reconstructed track.

Track selection

To have a clean sample of tracks for the efficiency calculation, we ask for a very pure sample of tracks, as summarised in table 2.6. Two sets of tracks are used: tracks reconstructed with combined information from the whole inner detectors (COMBINEDINDET tracks) and those from SCT measurements alone (SCTSTANDALONE tracks). Tracks in both sets are required to come from the interaction point via a cut on the impact parameter d_0 , defining the transverse distance to the primary vertex. Tracks must be fitted well by the pattern recognition algorithm via a cut on the χ^2 of the track, and are required to have $p_T > 1$ GeV. Also, an ‘X1X’ hit pattern, or level sensing, is required, meaning that only channels with a hit in the time-bin of the bunch crossing are read out. Finally, there should be at least 6 (7) other hits on the COMBINEDINDET (SCTSTANDALONE) track excluding the hit or hole under consideration. This last requirement builds confidence that the track indeed leaves a hit at (nearly) all intersections with the modules.

When comparing the evolution of the hit efficiency throughout the run of data taking period with a high number of interactions per bunch crossing with that of a run with low $\langle\mu\rangle$ (figure 2.12 (a)), an interesting observation is made. As a function of the respective luminosity block, which is a certain portion of data taking⁵, the hit efficiency is observed to increase during a run in data taking periods with high $\langle\mu\rangle$. These high $\langle\mu\rangle$ runs have a high number of tracks per event. This is illustrated in figure 2.12 (b), where the track multiplicity is shown for events from a run with $\langle\mu\rangle \sim 25$, and for a run with $\langle\mu\rangle < 1$. Events with many tracks leave a large number of hits in the SCT detector, or a high occupancy. This occupancy ranges between 1 and 2% per SCT layer in 2012 data taking. The tracking algorithm, described in section 4.1, may lead to a bias in these events: hits might be ascribed to tracks incorrectly, leading to ‘fake’ inefficiencies. Figure 2.12 (c) shows the measured SCT hit efficiency as a function

⁵In 2012 one luminosity block corresponded to 1 minute worth of data

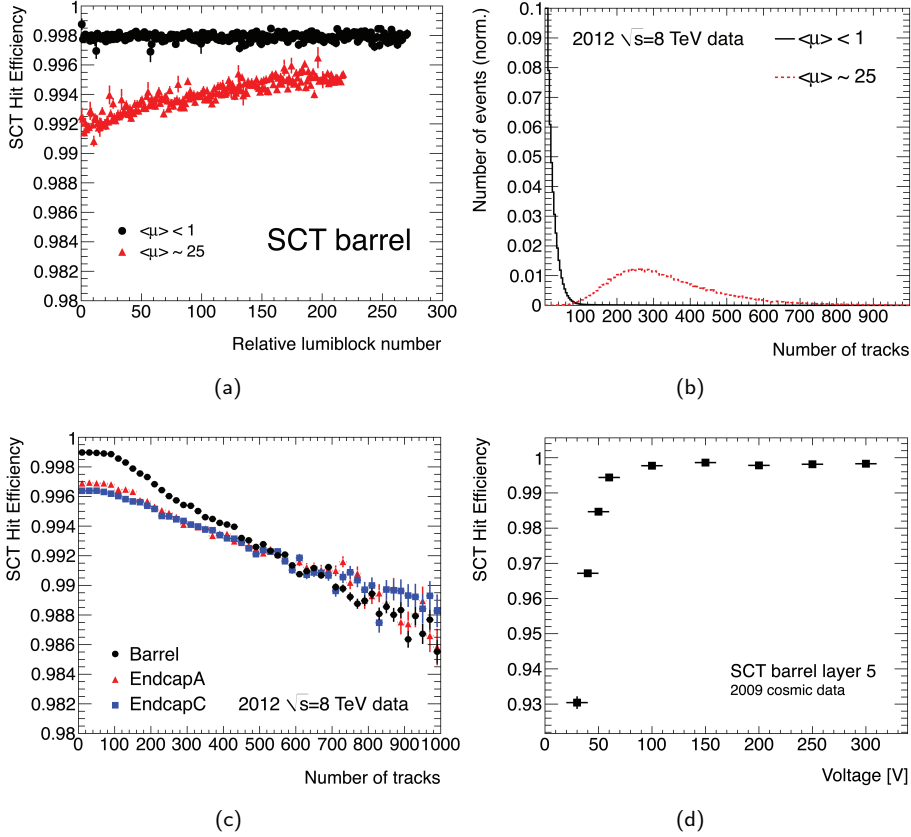


Figure 2.12: (a) Evolution of the SCT barrel hit efficiency throughout a run, as a function of the luminosity block (relative to the first with stable beams), measured in two runs: in red with a high mean number of interactions per collision $\langle\mu\rangle \sim 25$, taken from 2012 data at $\sqrt{s} = 8$ TeV; in black with $\langle\mu\rangle < 1$, measured in 2010 at $\sqrt{s} = 7$ TeV. (b) Measured number of tracks per event for the same data. (c) SCT hit efficiency as a function of the number of tracks per event, for barrel (circles), end-cap A (triangles) and end-cap C modules (squares), measured in a run with $\langle\mu\rangle < 1$. (d) SCT hit efficiency of the barrel as a function of the high voltage, measured in data taken from cosmic muons. Error bars show statistical uncertainties.

of the number of tracks. The efficiency decreases with increasing number of tracks, confirming the above assumption. For this reason, only events with less than 250 tracks are used for the hit efficiency calculation. During the run, the instantaneous luminosity decreases due to both proton-proton and proton-gas collisions, leading to a lower number of pp collisions per bunch crossing and thus less tracks, explaining the

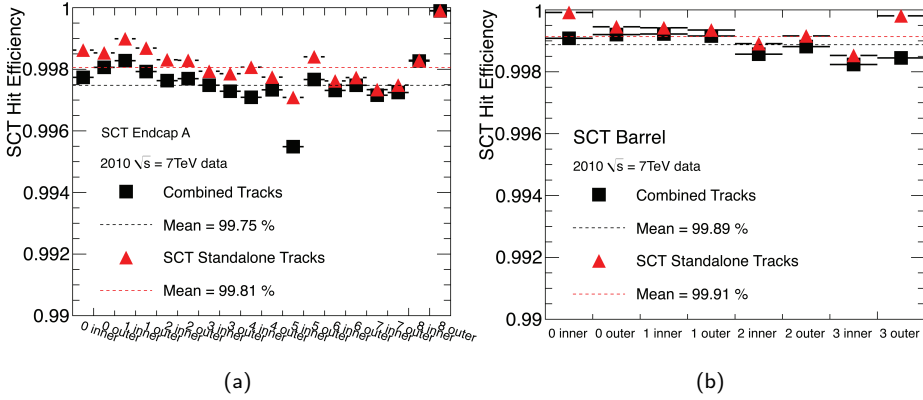


Figure 2.13: The SCT hit efficiency of (a) end-cap A and (b) the barrel in 2010 run ‘165591’ at $\sqrt{s} = 7$ TeV, for COMBINEDINDET tracks and SCT-STANDALONE tracks. The efficiencies are shown for each individual layer. Error bars, showing statistical uncertainties, are smaller than the markers. The given mean efficiency excludes the fully efficient layers. Although not shown, the efficiency for end-cap C is very similar to that of end-cap A.

increase of hit efficiency seen in figure 2.12 (a). To study the effect of the value of the bias voltage on the efficiency of the SCT, data from cosmic muons gathered in 2009 is used. During several of these cosmic data taking runs, the third layer of the barrel had a non-standard bias voltage applied. To be able to select cosmic muons travelling through the detector, the d_0 requirement is disregarded. Figure 2.12 (d) shows the hit efficiency of the barrel layer as a function of the bias voltage. A clear turn-on is observed: from 60 V onwards the efficiency is above 99%, while it seems to be on a plateau from 100 V, validating the decision to reject data taken during the warm start of the detector, where the bias voltage is still increasing but not yet at the standard 150 V.

During the run the instantaneous luminosity decreases due to both proton-proton and proton-gas collisions, leading to a lower number of pp collisions per bunch crossing and thus less tracks.

Hit efficiency during Run 1

The calculated intrinsic efficiency of the SCT changes slightly between runs, mainly due to changes in the configuration of the detector such as differences in disabled strips. Figure 2.13 shows the hit efficiency for a typical run in 2010 with $\sqrt{s} = 7$ TeV data for the barrel and both end-caps, comparing the efficiency with COMBINEDINDET tracks to SCTSTANDALONE tracks. The mean efficiency is above 99.7% for all components, while the SCTSTANDALONE tracks have a somewhat higher efficiency due to the tighter cut on the number of hits. Looking closely, layers at the end and

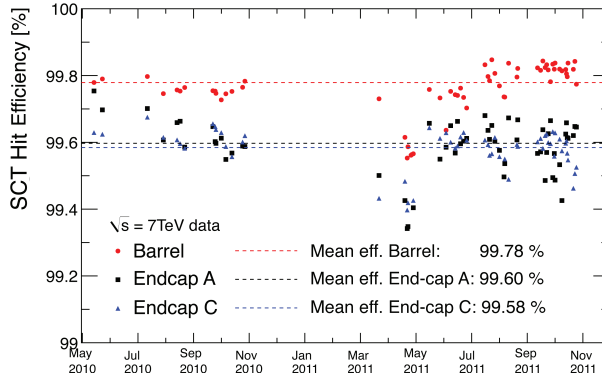


Figure 2.14: Trend of the SCT hit efficiency during runs in 2010 and 2011 at $\sqrt{s} = 7$ TeV, for the barrel (red), end-cap A (black) and end-cap C (blue). The mean values are given by the dashed lines. Error bars, showing statistical uncertainties, are smaller than the markers.

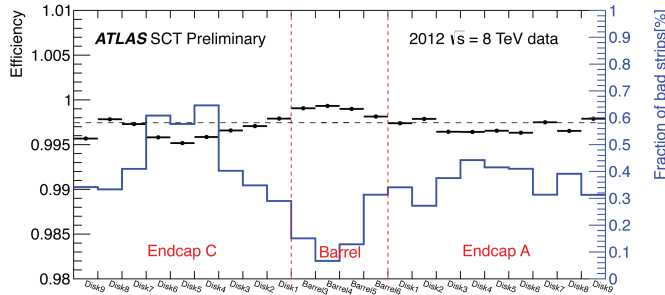


Figure 2.15: SCT hit efficiency (markers) in a typical run in 2012 at $\sqrt{s} = 8$ TeV, together with the fraction of disabled strips (histogram) [138]. Error bars, showing statistical uncertainties, are smaller than the markers.

beginning of tracks are fully efficient: for COMBINEDINDET tracks these are just the outer end-cap layers, while for SCTSTANDALONE tracks the inner and outer layers of both barrel and end-cap can be the beginning or end of a track. These layers artificially obtain a 100% efficiency due to the fact that the efficiency calculation interprets the first (last) hit on track to be the (beginning) end of the track itself; any hole before or after the track is not counted as a hole, while a hit *is* counted in the calculation. For this reason, the mean efficiency given in the figures excludes these layers.

To investigate the stability of the SCT during the first two years of LHC operation, and thus get more than a snapshot of the SCT performance, the trend in the efficiency is plotted in figure 2.14. This trend is obtained by calculating the average efficiency of the barrel (shown in red), end-cap A (black) and end-cap C (blue) for many runs

under the same conditions. For each data taking sub-period⁶ two runs with the largest number of events are selected. The intrinsic efficiency is seen to be stable over the two first years of running for each of the components. A small dip can be seen at the start of the 2011 data taking, coming from a combination of the fact that the SCT was recording data in 'XXX' mode, using data with hits outside the bunch-crossing time-bin, and a higher number of disabled strips. In May the SCT was reconfigured to 'X1X' level sensing. Alignment improvements introduced July 2011 lead to the increase in efficiency seen in the last few months.

Finally, the SCT hit efficiency in 2012 has not deteriorated. The intrinsic efficiency for each SCT layer is shown in figure 2.15 for a typical run in 2012. The mean value, represented by the dashed line, is $99.74 \pm 0.04\%$, where the error is a systematic uncertainty derived from varying the distance cut for non-associated hits and track selection criteria. Again, the efficiency for barrel modules is about 0.2% higher than for end-cap modules. This is explained by studying the fraction of inactive strips, shown in the blue histogram. Modules in both end-caps have a higher number of isolated disabled strips than modules in the barrel, and even a anti-correlation between efficiency and number of inactive strips is seen. Thus differences in efficiency between layers seem to predominantly be due to relative differences in the number of inactive strips.

2.8.2 Operational status of the other sub-detectors

The various sub-detectors have mostly been performing well. Figure 2.16 (a) shows the hit efficiency of each of the layers of the pixel detector for data taken in 2010, with $\sqrt{s} = 7$ TeV. This gives the probability of finding a hit in a pixel layer after a track has crossed it. It can be seen that all layers except the two outer end-caps have an efficiency around 99%. The outer layers are affected by some dead or inefficient regions on some modules, but are still above 97% efficient. The full efficiency of the inner B-layer is due to the track selection, and thus is 100% by construction. The detector noise is under control, as can be seen by the noise occupancy shown in figure 2.16 (b). The noise occupancy gives the probability of detecting a random hit not associated with a track. It is obtained by measuring the detector occupancy in empty bunch crossings. After reconstruction there are less than 10^{-9} noise hits per pixel per bunch crossing, or less than 0.04 hits per event.

The hit efficiency of the barrel of the TRT is shown in figure 2.16 (c) as a function of the track's distance of closest approach to the wire. The hit efficiency is defined as the number of straws traversed by the track with a hit, divided by the total number of straws traversed by the track. Here the known disabled straws are excluded from the calculation. The mean of the efficiency is above 94% in the plateau region – only for tracks far away from any straw centre the efficiency drops. On the right,

⁶For every significant change in detector hardware or software, or in LHC conditions, a new data taking period is defined. These are labelled by letters, starting from period A. The data taking periods are again subdivided into smaller sub-periods A1, A2, ... etc. for minor changes in the conditions.

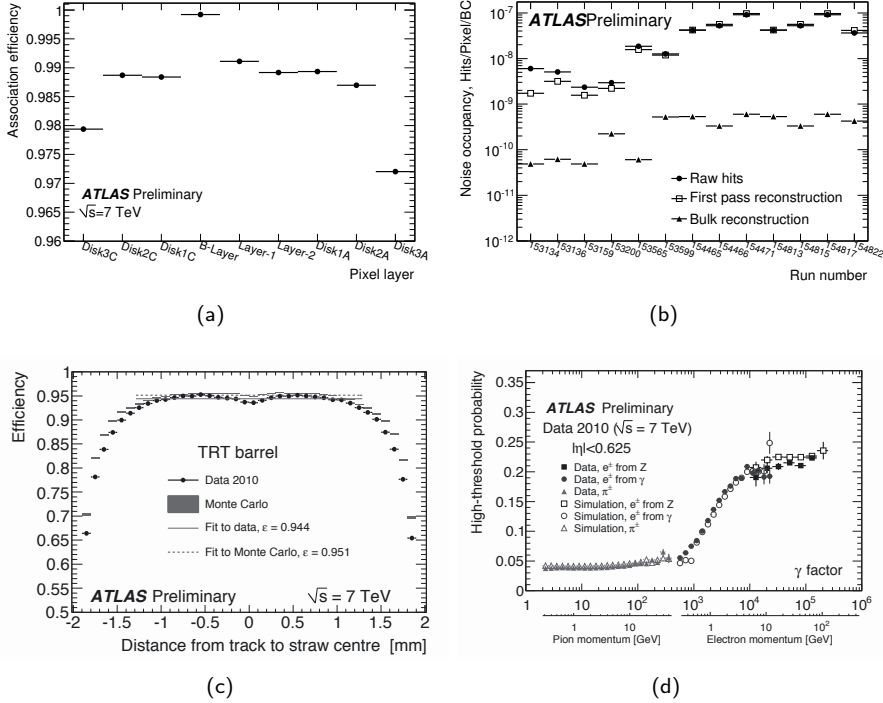


Figure 2.16: (a) Measured hit efficiency of the pixel detector. The non-functioning pixels are excluded from the efficiency calculations. (b) Measured noise occupancy of the pixel detector. After reconstruction the noise occupancy is $< 10^{-9}$ hits per pixel per bunch crossing. (c) Hit efficiency of the TRT barrel as a function of the distance of the track to the straw centre, with a mean of 94.4% in the plateau region. Data taken in 2010 is compared to Monte Carlo simulations. The non-functioning straws are excluded in the efficiency calculation. (d) Probability for a transition radiation hit producing a signal over the threshold in the TRT barrel as a function of their momentum and γ factor. Tracks from pions are shown left (triangles), electrons coming from Z boson decays (black) and photon conversions (blue) are shown right. Data is compared to Monte Carlo simulations (open markers). All data taken in 2010 at $\sqrt{s} = 7$ TeV. Figures taken from [139].

in figure 2.16 (d), the probability for a track creating transition radiation to produce a signal above the high-threshold is shown, for both pion tracks (left) and electron tracks (right). It is given as a function of track momentum, or alternatively γ factor. The data used was taken with $\sqrt{s} = 7$ TeV data in 2010, and is compared to a minimum bias Monte Carlo simulation. Electrons have up to 25% chance of giving a high-threshold hit in a given straw tube, of which they traverse up to 34 per track,

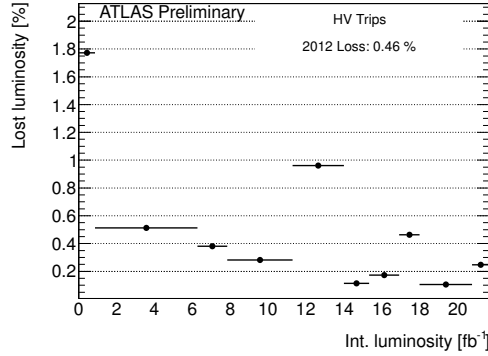


Figure 2.17: The data loss in the various data taking periods due to HV trips in the liquid argon calorimeters in 2012.

enabling the discrimination of electrons and pions.

The liquid argon calorimeters, ECAL, HEC and FCal, all have more than 99.5% of the channels operational. The main source of loss of data in these calorimeters is due to high voltage trips. This is shown in figure 2.17 as a function of the gathered integrated luminosity in 2012. Apart from the very first period, which only consisted of a small fraction of the total gathered data, none of the data taking periods suffered a loss of more than 1% of data due to these trips, with an average of 0.46% over the whole year. The total loss of data, including other sources, was 0.88%. In May and June of 2011, 6 front-end boards (FEBs) were non-operational due to an electronics failure. This affected the second and third layer of the ECAL in the region $0 < \eta < 1.4$ and $-0.8 < \phi < -0.6$, where no energy could be measured. From July onwards, 4 out of the 6 FEBs were recovered, recovering the full second layer. As shown in table 2.4, on average about 98% of the number of cells has been operational in the Tile calorimeter. Yet in between shut-downs, without access to the detector, the low voltage power supply failed for several units containing front-end electronics, leading to a loss of operational cells of up to 5% of total at the end of 2011.

Without counting the so-called EE chambers (see figure 2.10) which are not yet installed, the fraction of operational channels in the muon spectrometer is above 96%. The uncertainty on the momentum resolution from the MS alignment is somewhat higher than the design level given in table 2.2 at $\sigma(p_T)/p_T = 13\%$ TeV^{-1} for the barrel, $\sigma(p_T)/p_T = 17\%$ TeV^{-1} for the MDT end-caps and $\sigma(p_T)/p_T = 14\%$ TeV^{-1} for the CSC end-caps [140].

2.8.3 Pile-up and vertexing in ATLAS

As mentioned in section 2.1.1, the conditions for the ATLAS detector have changed during the three years of data taking. Arguably the most important change has been the evolution of the average number of collisions per bunch crossing, which has in-

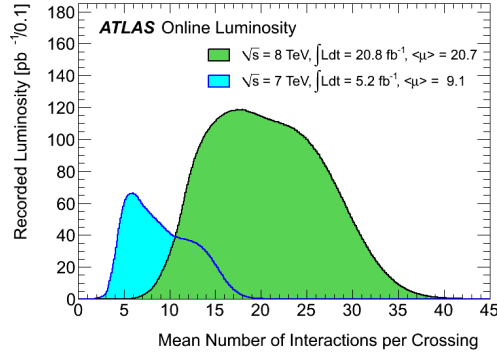


Figure 2.18: The amount of recorded integrated luminosity as a function of the mean number of interactions per bunch crossing for collisions in the ATLAS detector in 2011 (blue) and 2012 (green). The average over 2011 was 9.1 interactions per crossing, while over 2012 this was 20.7.

creased dramatically from around one in 2010 to more than 20 at the end of 2012. Figure 2.18 shows the amount of recorded integrated luminosity as a function of the mean number of interactions per bunch crossing $\langle \mu \rangle$, for both 2011 and 2012. The rise is clearly visible, with an average $\langle \mu \rangle = 20.7$ over 2012. Here μ is calculated in data as [141]

$$\mu = \frac{L \times \sigma_{\text{inel}}}{n_{f_r, \text{bunch}}}, \quad (2.2)$$

with L the integrated luminosity, σ_{inel} the total inelastic cross section, n_{bunch} the number of colliding bunches and f_r the revolution frequency of the protons. The inelastic cross section is taken as 71.5 mb from PYTHIA [142], with a 3% uncertainty due to the differences between this value and the measured value of 73.5 ± 1.9 mb by TOTEM [143] and $69.1 \pm 2.4 \pm 6.9$ mb by ATLAS [144]. The uncertainties on the ATLAS measurement are the experimental and extrapolation uncertainties, respectively.

The high mean number of collisions per bunch crossing has several implications for physics analyses. As a positive effect, and the main reason for this configuration, a higher collision rate means a higher number of interesting events. Yet this comes at a price. First of all, the trigger encounters more interesting events, and due to the restrictions on the output rate more stringent thresholds need to be set. More interactions also lead to more charged tracks in the inner detector, which need to be reconstructed, and more energy deposits in the calorimeters. By reconstructing the primary vertex (the vertex of interest) with high precision, and tracking the particles on their path to the calorimeters, the separate events can be disentangled. The pile-up causes jets to have extra energy, while the $\cancel{E}_T$ resolution decreases. The effects of this in-time pile-up on the jet reconstruction are discussed further in section 4.2.

SUSY Cross Sections

Assuming that SUSY with R-parity conservation is realised in nature, SUSY-particles will be pair produced at the LHC if their masses are not too large. The production of coloured sparticles, the gluinos ($\tilde{g}$) and squarks ($\tilde{q}$), are expected to be the main production channel:

$$pp \rightarrow \tilde{g}\tilde{g}, \tilde{q}\tilde{q}, \tilde{q}\tilde{\bar{q}}, \tilde{\bar{q}}\tilde{q}, \quad (3.1)$$

whereas the electroweak sparticle production will happen less often.

The fraction of LHC collisions producing a SUSY-particle pair $\tilde{X}\tilde{Y}$ is determined by the cross section $\sigma_{\tilde{X}\tilde{Y}}$ of the process $pp \rightarrow \tilde{X}\tilde{Y}(+anything)$. An accurate theoretical prediction of these cross sections is imperative in the search for SUSY at the LHC (and other colliders) for the interpretation of results in a SUSY-framework (see chapter 5). This chapter describes the method used within ATLAS to calculate these cross sections and the uncertainties on them. The first section discusses the motivation and theoretical framework; the second section goes into details on the various contributions of the theoretical uncertainty on the cross section, while the third section finally combines these into two different methods of calculating the final cross section and its uncertainty, and goes into details for several specific SUSY models used in ATLAS analyses.

3.1 Theoretical introduction

As discussed in section 1.1.2, the colliding protons at the LHC are composite particles consisting of 3 valence quarks, embedded in a sea of virtual quark pairs, with the gluons holding them together – making QCD the underlying theory for LHC collisions. Although we measure a hadronic cross section (for Standard Model processes), we can only calculate the underlying partonic cross sections in perturbative QCD. A translation has to be made between these to have any predictive power over a measurement. This translation is done by factorising long- and short-distance effects from each other, where the short-distance effects are given by the partonic cross section and the long-distance effects are included in parton distribution functions (PDFs). This factorisation is shown in equation 1.29.

The quality of the prediction of the cross section comes from how well the partonic cross section can be calculated, and how well the PDFs are known. These PDFs have been derived by various collaborations, mainly CTEQ [33], MSTW [32] and NNPDF [34], from recent measurements up to a certain accuracy. The remaining uncertainty in the PDFs causes a systematic uncertainty in $\sigma_{\tilde{X}\tilde{Y}}$, studied in section 3.3.1. The uncertainty on the α_S measurements leads to a minor contribution, which is shown in section 3.3.2.

The factorisation and renormalisation scales, μ_F and μ_R , are usually taken to be both equal to the typical energy scale of the process at hand, i.e. the average mass of the outgoing partons: $\mu = \mu_F = \mu_R = Q$, with e.g. $Q = m_{\tilde{q}}$ for squark-pair production. The dependence of $\sigma_{\tilde{X}\tilde{Y}}$ on μ results in an additional theoretical uncertainty (see section 3.3.3). The uncertainty due to the dependence on the masses of the Standard Model particles, foremost the top quark mass m_t , is not studied here.

3.1.1 Next-to-leading order and beyond

As stated before, the hard part of the cross section can in principle be calculated to all orders of α_S in perturbation theory, yet in practice the limitations show up quickly. Confinement prescribes that the QCD coupling constant α_S will be large for small energy scales, thus leading perturbation theory to fail for small energy scales. For low energy QCD, lattice calculations are the only possibility. For the high energy scales at the LHC α_S is sufficiently low, yet this does not remove all obstacles. For each next order of perturbation theory, the number of diagrams which will have to be calculated increases dramatically, making calculations beyond leading order (LO) tedious, while every order beyond that will be even more difficult. For the SUSY-QCD processes discussed above they have been available in next-to-leading order (NLO) for some time [83, 145–147]. Adding higher-order terms reduces the scale dependence of the cross section, as will be discussed in section 3.3.3.

A large part of the NLO corrections to the cross section comes from the so-called threshold region, where the partonic centre-of-mass energy is close to the kinematic threshold of the production. Consider the production of two heavy particles. The LO cross section of this production is zero for energies just below the mass of the system, as one cannot produce the system without enough energy. Increasing the energy above this threshold, the cross section rises steeply at first, giving rise to large differences in the cross section for small energy changes. When going to next-to-leading order, diagrams with radiation of particles are included. These radiation effects depend on the energy of the radiated particle: the probability to radiate a high-energy particle is small, while low-energy particles are radiated easily. The soft-gluon radiation off coloured particles will affect the energy available for the production of the heavy particles, making this radiation very important at the threshold region. We expect the mass of SUSY particles to be a significant fraction of the total collision energy, and the stochastic nature of the distribution of the energy of the proton over its constituent partons assures that the probability of producing a SUSY particle near the threshold region is large – that is, if such a SUSY particle exists. This threshold region

is therefore very important for new physics searches at the LHC [148].

The emission of soft-gluons off the initial and final state gives rise to higher-order QCD corrections of the form [149]:

$$\alpha_S^n \log^m \beta^2, m \leq 2n \text{ with } \beta^2 = 1 - \frac{(m_{\tilde{X}} + m_{\tilde{Y}})^2}{s}, \quad (3.2)$$

where $m_{\tilde{X}}$ and $m_{\tilde{Y}}$ are the masses of the two final state particles, s is the partonic centre-of-mass energy squared, and the integer n gives the order of perturbation in QCD, while m gives the order in the logarithmic perturbation. The partonic centre-of-mass energy is found by using $s = x_1 x_2 S$, where x_1, x_2 are the parton momentum fractions of partons 1 and 2, and S is the hadronic centre-of-mass energy squared. From the above it is obvious that once s approaches the masses of the final states, and thus $\beta \rightarrow 0$, these corrections can become very large, and the perturbative expansion will have convergence issues.

These issues can be solved by using the threshold resummation technique demonstrated in ref. [149–152]. Without going into theoretical detail, this technique takes the soft-gluon emission into account for all orders in α_S by using a new perturbative expansion in the large logarithms. For gluon-pair and squark-squark production the state-of-the-art calculations go up to next-to-leading logarithmic order (NLL), while for squark-antisquark production next-to-next-to-leading logarithmic order (NNLL) is even available [153], although in this thesis only the NLL terms are used.

This results in two parts of the calculated partonic cross section: the next-to-leading order terms of the hadronic cross section, and the next-to-leading logarithmic terms for the soft-gluon resummation. To combine these, a matching procedure has to be applied, see [149], resulting in a NLO+NLL cross section calculation. In this chapter soft-gluon resummation is taken into account only where stated, for instance in the advanced procedure for obtaining cross section for ATLAS analyses in section 3.4.2. In the following sections, the two programs which are used to calculate NLO and NLO+NLL cross sections are introduced: PROSPINO and NLL-FAST respectively.

3.1.2 Prospino

To obtain the next-to-leading order cross sections for SUSY sparticle production, we use PROSPINO2.1 [74, 83, 147, 154–156], which analytically calculates the partonic cross section to next-to-leading order in the strong coupling constant, α_S . It includes calculations for the most abundant production processes at the LHC, namely the pair production of squarks, gluinos, top squarks, gauginos and sleptons. Furthermore the associated squark-gluino production and the production of a gaugino with a squark or a gluino is included. PROSPINO can be interfaced with NLO PDFs from the CTEQ and MSTW collaborations to obtain hadronic cross sections; the NNPDF PDFs have not yet been included.

In equation 3.1 the squark chiralities have been suppressed for simplicity, $\tilde{q} = (\tilde{q}_L, \tilde{q}_R)$. More importantly, it hides the three squark generations. In PROSPINO the stop, the supersymmetric partner of the top, is treated separately from the part-

ners of the five lightest quarks ($\tilde{u}, \tilde{d}, \tilde{c}, \tilde{s}, \tilde{b}$) due to large mixing effects and large mass differences between the mass eigenstates $\tilde{t}_1, \tilde{t}_2$. Also the negligible fraction of top quarks in the proton and thus in the PDFs removes corresponding diagrams from the calculation. Although the same holds to lesser extent for the sbottoms, they are usually treated as light squarks since the detector has difficulty in separating light-flavour jets from bottom decay jets.

In principle we calculate a mass degenerate squark-(anti)squark cross section for the five lightest flavour left- and right-handed squarks, i.e. one summed cross section where every squark has a mass equal to the average of all five lightest flavour left- and right-handed squarks. To calculate the individual squark-pair cross sections in PROSPINO, the NLO cross sections for the five light flavour squarks are calculated using an assumption on the K-factors, i.e. the ratio of NLO to LO cross sections. First, their individual exact LO cross sections are calculated, which are multiplied by one mass degenerate K factor. This mass degenerate K-factor is derived by dividing NLO by LO cross section, both with all five light flavour squarks set to their average mass. Assuming this K-factor is equal to the separate K-factors for each individual process, we obtain the NLO cross section for the individual squarks.

3.1.3 NLL-Fast

For the resummation of the soft gluon emission into the cross section calculation, the numerical code NLL-FAST has been written [157, 158]¹ which adds the soft-gluon resummation to the next-to-leading order calculations performed by PROSPINO to obtain NLO+NLL cross sections for the production of the coloured SUSY particles. For computational efficiency the central value cross sections have been arranged in tables for pairs of coloured sparticles as a function of their masses. The same is done for the scale uncertainties, PDF uncertainties at 68% CL and strong coupling uncertainties. These tables can be used to look up central value cross sections and their uncertainty directly for a given sparticle mass. However, the granularity of the table is chosen such that one can also interpolate with high accuracy ($< 1\%$) between the mass values given in the table. The usage of NLO PDFs in the NLL calculation is supported by the argument that the evolution of the two have the same divergences [148, 160]. In this thesis, wherever NLO+NLL cross sections are mentioned this package is used.

3.2 NLO SUSY cross sections

Using PROSPINO we can calculate the NLO cross sections for a grid of points in the CMSSM for the various production processes allowed. Figure 3.1 shows the cross section for a CMSSM model where 3 of the 5 parameters are fixed: $\tan\beta = 10$, $A_0 = 0$ and the sign of μ is positive. The remaining two parameters, the common scalar mass m_0 and common gaugino mass $m_{1/2}$, define the 2-dimensional plane shown. This model is used in both the 2011 analysis performed on 4.7 fb⁻¹ of $\sqrt{s} = 7$ TeV data, as well as the $\sqrt{s} = 8$ TeV analysis described in chapter 5 in this thesis. For the

¹Available from ref. [159]

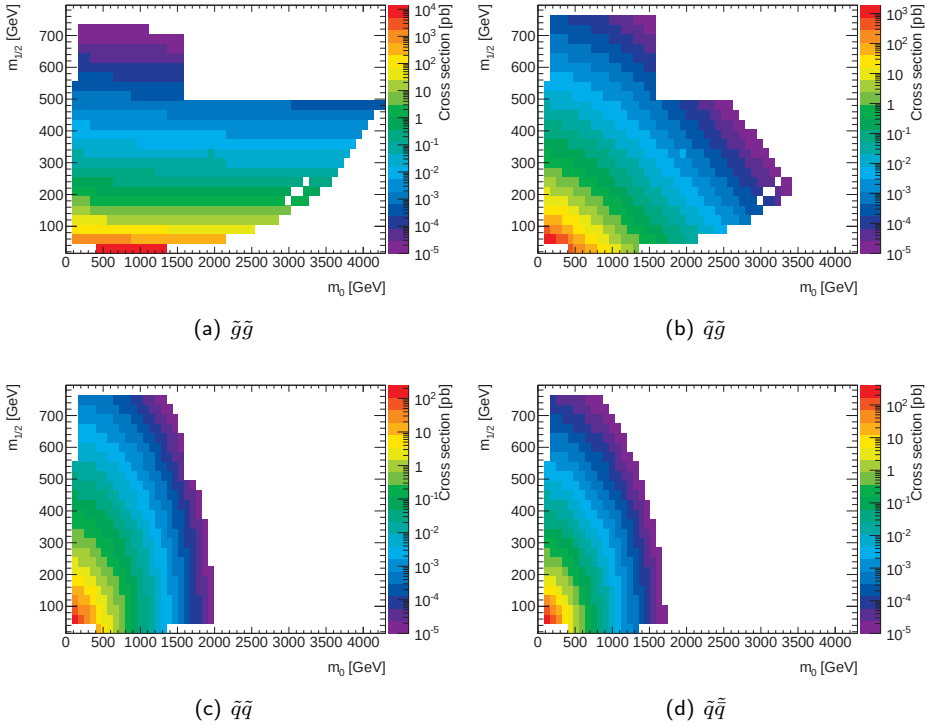


Figure 3.1: The next-to-leading order cross sections in picobarn calculated with PROSPINO using CTEQ6.6 PDFs (the ‘straightforward’ method) for various production processes in the CMSSM plane at $\sqrt{s} = 7$ TeV: (a) gluino-pair production, (b) associated squark-gluino production, (c) squark-pair production and (d) squark-antisquark production. Note that the scale on the z -axis goes to higher values for gluino-pair and gluino-squark production. The cross section is only shown down to 10^{-5} pb.

calculation of the values of the cross sections the central value CTEQ6.6 PDFs are used, as these PDFs are used for the generation of simulated events in ATLAS.

Although the CMSSM allows for the production of the complete spectrum of SUSY particles (unlike the simplified models), figure 3.1 shows the results for $\tilde{g}\tilde{g}$, $\tilde{q}\tilde{g}$, $\tilde{q}\tilde{q}$ and $\tilde{q}\tilde{\bar{q}}$ production. Note that as mentioned before, the squarks include the five lightest flavours. These four processes have on average the highest cross section in the CMSSM, although light electroweak SUSY particles can have a significant contribution to the total SUSY cross section for low $m_{1/2}$.

The cross section is seen to depend heavily on the mass of the produced sparticle, as is expected. Furthermore, it becomes clear that gluino-pair production and associated gluino-squark production contribute most significantly across large part of the phase

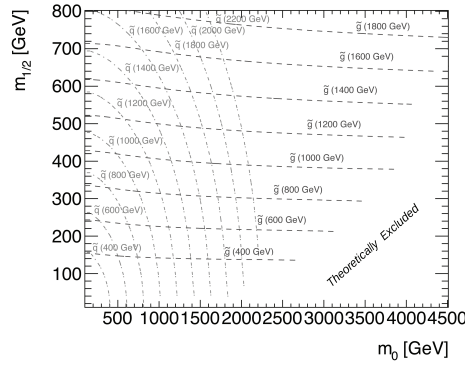


Figure 3.2: *Dependency of the squark and gluino masses on m_0 and $m_{1/2}$ in an CMSSM model with $\tan\beta = 10$, $A_0 = 0$ and μ is positive. The region on the bottom right is theoretically excluded.*

space, with the squark-pair production only being dominant at high $m_{1/2}$ and very low m_0 , which translates into the highest squark-gluino mass ratio. The squark-antisquark production cross section has the same dependence on m_0 and $m_{1/2}$ as that of the squark-pair production, yet the value of the cross section is about one order of magnitude less. The dependency of the cross section on the masses of the gluinos and squark becomes evident when one compares the shapes in these plots with the m_0 , $m_{1/2}$ dependency of the squark and gluino masses of the particles shown in figure 3.2. Note that for the most contributing process, $\tilde{q}\tilde{q}$, the cross section decreases by 8 orders of magnitude at high $m_{1/2}$. This is a common feature for all processes, and is one of the main reasons analyses need a vast amount statistics to gain more sensitivity at high masses.

3.3 Contributions to the cross section uncertainties

As discussed in the introduction, the calculation of the cross section for the production of supersymmetric particles depends on several variables, which need to be measured in separate experiments. For instance, the strong coupling constant α_S has been measured in many experiments, like deep inelastic scattering processes, hadronic τ decays, electron-positron annihilation processes and electro-weak precision fits [161]. The theoretical uncertainty on the SUSY cross sections are due to three sources:

- The parton distribution functions
- The renormalisation and factorisation scale μ
- The strong coupling constant.

In this section these contributions to the uncertainty on the cross section are discussed.

3.3.1 PDF variations

For the calculation of next-to-leading order (NLO) cross sections, NLO PDFs are used, from either CTEQ6.6 [33] or MSTW2008 [32]. These PDF sets are extrapolated from measurements by the corresponding collaborations (section 1.1.2), and thus carry uncertainties. The PDF fits done by the collaborations lead to 22 (20) eigenvectors spanning the range of uncertainties coming from the experimental errors at 90% CL (68% CL) for the CTEQ (MSTW) collaboration. The PDF sets contain PDFs for the gluon and lightest 5 quarks and anti-quarks. The collaborations supply for each of these a central value PDF (corresponding to the best fit), together with an upward or downward varied PDF for each of its eigenvectors. This results in 44 and 40 PDFs for the CTEQ and MSTW collaborations respectively. We can now use the Hessian method [42] to obtain a symmetric PDF uncertainty on our cross section. As the cross section depends on the PDFs, we can characterise the cross section by a function $\sigma(\vec{a})$, where $\vec{a}$ is a vector in the PDF-parameter space. We can now approximate the change of the cross section along $\vec{a}$ as

$$\frac{\partial \sigma}{\partial a_i} = \frac{1}{2}(\sigma_i^+ - \sigma_i^-) \quad (3.3)$$

where $\sigma_i^\pm$ are the upper and lower result of the cross section for the i th eigenvector of the PDF set. To obtain the variation of $\sigma(\vec{a})$ over all eigenvectors, we can just add them quadratically [33, 42]:

$$\delta_{PDF}\sigma = \frac{1}{2}\sqrt{\sum_i (\sigma_i^+ - \sigma_i^-)^2}. \quad (3.4)$$

This means that to quantify how these uncertainties are propagated into the cross section of a specific sub-process, each of the PDF parameters is varied during the calculation of $\sigma_{\tilde{X}\tilde{Y}}$ while keeping the other parameters fixed, where $\tilde{X}\tilde{Y}$ is the produced sparticle pair. This leads to 44 (40) different outcomes of the cross section $\sigma_i^\pm$ for CTEQ (MSTW) PDFs. The resulting systematic uncertainty $\delta_{PDF, \tilde{X}\tilde{Y}}$ is then:

$$\delta_{PDF, \tilde{X}\tilde{Y}} = \frac{1}{2}\sqrt{\sum_i (\sigma_i^+ - \sigma_i^-)^2} \quad (3.5)$$

for each sub-process separately. The relative uncertainty for channel $\tilde{X}\tilde{Y}$ is just $\Delta_{PDF, \tilde{X}\tilde{Y}} = \frac{\delta_{PDF, \tilde{X}\tilde{Y}}}{\sigma_{\tilde{X}\tilde{Y}}} \times 100\%$. In contrast, up- and downward PDF uncertainties can be constructed by only adding the up and down variations quadratically, respectively:

$$\delta_{PDF}^+ \sigma = \sqrt{\sum_i (\sigma_i^+ - \sigma_i)^2} \quad (3.6)$$

$$\delta_{PDF}^- \sigma = \sqrt{\sum_i (\sigma_i - \sigma_i^-)^2}. \quad (3.7)$$

These asymmetric uncertainties will be more conservative than those obtained by the

Hessian method. As mentioned, for the CTEQ PDFs this is still an uncertainty at a 90% CL. To obtain the uncertainty at 68% CL (or one standard deviation) $\Delta_{PDF, \tilde{X}\tilde{Y}}$ has to be rescaled with 1.645.

In figure 3.3 the PDF uncertainty Δ_{PDF} is shown for $pp \rightarrow \tilde{q}\tilde{g}, \tilde{g}\tilde{g}, \tilde{q}\tilde{q}$ and $\tilde{q}\tilde{\bar{q}}$. The results are shown in the CMSSM plane. In these, large differences are seen across the plane for each process, and also between the different processes. Focussing on the variations over the grid of points for the shown processes, we see that for gluino-pair production (figure 3.3 (a)) the uncertainty grows approximately linearly with $m_{1/2}$, or more accurately, with the gluino mass. The same is seen for the squark-pair production, where the PDF uncertainty increases with increasing $m_{\tilde{q}}$. This can be understood when considering how PDFs are obtained: they are measured from experimental data at certain energy scales and for ranges in momentum fraction $10^{-4} < x < 0.1$. Beyond these energies and x ranges, they are extrapolated using the DGLAP evolution, yet the further from the data, the larger the uncertainties become.

The difference between the various processes is also very apparent. Where for the gluino-pair production the PDF uncertainties range between 10-70%, the squark-pair production has a smaller PDF uncertainty of 20% or below for the most important region. This difference is a feature of the production of sparticles from protons: gluinos are mostly produced by gluon fusion, while the squarks are produced from interactions between the valence quarks within the proton. The uncertainty on gluon PDFs is larger than the one on the quark PDFs for low values of x , leading to higher uncertainties on the cross section.

Note that the shown uncertainty is the symmetric PDF uncertainty. Using equations 3.6 and 3.7 the separate up- and downward uncertainties can be calculated, which can be smaller than the symmetric uncertainty.

3.3.2 Variations in the strong coupling

The world average value of α_S based on eight different measurements is $\alpha_S(M_Z^2) = 0.1184 \pm 0.0007$ [161], evaluated at the mass of the Z-boson. The various PDF collaborations use different values: CTEQ and NNPDF adopt the world average, while MSTW actually fits the value of $\alpha_S(m_Z^2)$ simultaneously with their fit of the PDF parameters, and then use the best fit value: $\alpha_S(M_Z^2) = 0.1207$. The uncertainty on the used value of $\alpha_S(M_Z^2)$ propagates into a second uncertainty on the calculated cross section for SUSY sparticle production. Note that in ATLAS the uncertainty due to α_S is currently only used for the CTEQ6.6 PDFs due to implementation difficulties with the MSTW PDF sets.

Although PDFs are dependent on the value of α_S , the uncertainty Δ_{α_S} due to the strong coupling constant can be straightforwardly accounted for with a simple method, reproducing the correlation between the PDFs and α_S correctly [34, 162]. By using two special PDF sets AS_{-2} and AS_{+2} where the value of the coupling constant is set to $\alpha_S = 0.116$ and 0.120 respectively we can get the deviation in the cross section with a Confidence Level (CL) of 90%. These values correspond to (slightly more) than two standard deviations from the world average. Adding this deviation in quadrature

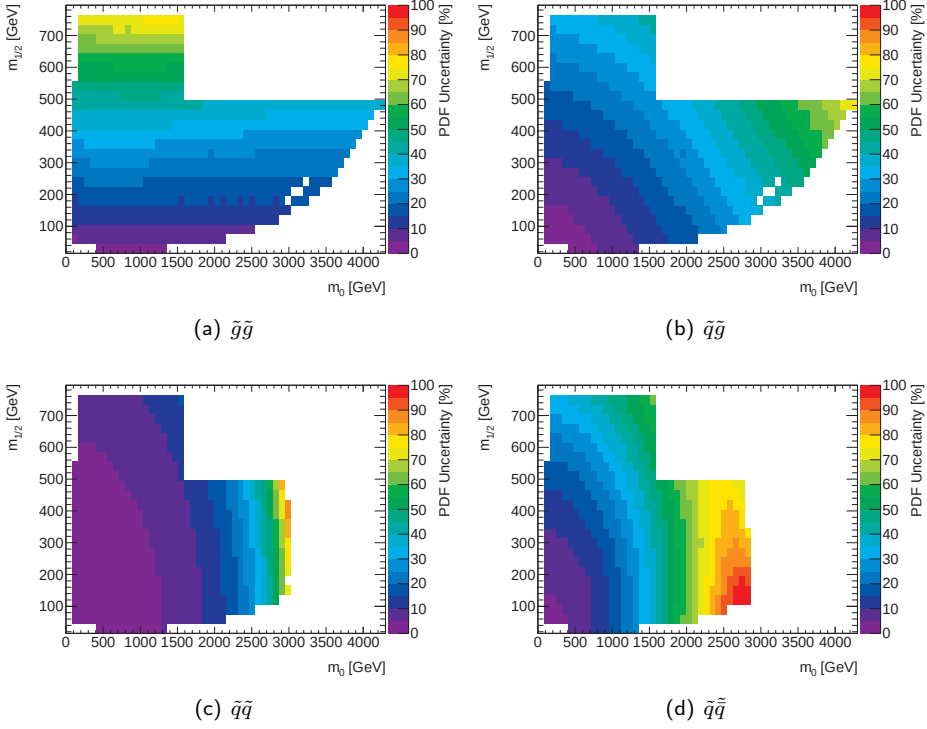


Figure 3.3: The relative PDF uncertainty $\Delta_{PDF, CTEQ}$ on next-to-leading order cross sections calculated with PROSPINO using CTEQ6.6 PDFs in%. Shown are various production processes in the CMSSM plane at $\sqrt{s} = 7$ TeV: (a) gluino-pair production, (b) associated squark-gluino production, (c) squark-pair production and (d) squark-antisquark production. For $\tilde{q}\tilde{q}$ and $\tilde{q}\tilde{q}$ production the uncertainties are shown until $m_0 \sim 3000$ GeV due to the low cross sections there.

to the uncertainty coming from the PDF variations (where the strong coupling is taken constant) will give a realistic determination of their combined uncertainty, as seen in ref. [162].

For the CMSSM plane, the uncertainty on the strong coupling constant is much smaller than that on the PDFs. Figure 3.4 shows the upward uncertainty for $\tilde{q}\tilde{g}$, $\tilde{g}\tilde{g}$, $\tilde{q}\tilde{q}$ and $\tilde{q}\tilde{q}$ production. The downward uncertainty is similar.

3.3.3 Scale variations

The final variables on which the calculation of the theoretical prediction of the SUSY cross section depends are the renormalisation and factorisation scales μ_R and μ_F . Although the choice of the renormalisation scale is arbitrary, which would give a scale

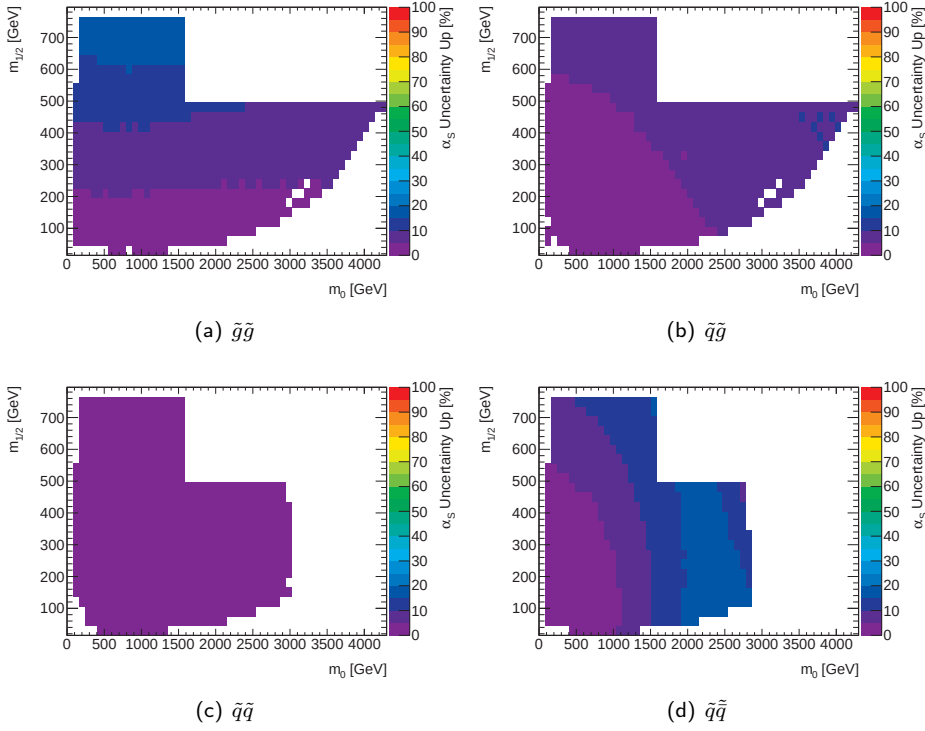


Figure 3.4: The relative upward α_S uncertainty Δ_{α_S} on next-to-leading order cross sections calculated with PROSPINO using CTEQ6.6 PDFs in%. Shown are various production processes in the CMSSM plane at $\sqrt{s} = 7$ TeV: (a) gluino-pair production, (b) associated squark-gluino production, (c) squark-pair production and (d) squark-antisquark production.

independent cross section, the cross section can only be calculated up to a certain order, leading to some scale dependence. We take μ_F and μ_R to be both equal² to the typical energy scale Q of the process at hand, $\mu = \mu_F = \mu_R = Q$. Here Q is taken to be the average mass of the final state sparticles, as the cross section is expected to depend mainly on this mass [163]: $Q = m_{\tilde{q}}$ for squark-(anti)squark production, $Q = m_{\tilde{g}}$ for gluino-pair production, $Q = \frac{1}{2}(m_{\tilde{q}} + m_{\tilde{g}})$ for associated squark-gluino production, and likewise for weak production.

The scale dependence at a certain order of perturbation theory should ideally contain the true value of the cross section, i.e. the cross section calculated to all orders of perturbation theory. Varying μ thus gives an estimation of the difference between σ^{NLO} and σ^{full} . Although the values between which μ should be varied are arbitrary,

²Varying μ_F and μ_R independently is outside the scope of this thesis. It is one of the priorities of the joint effort between ATLAS, CMS and the LPCC SUSY working group.

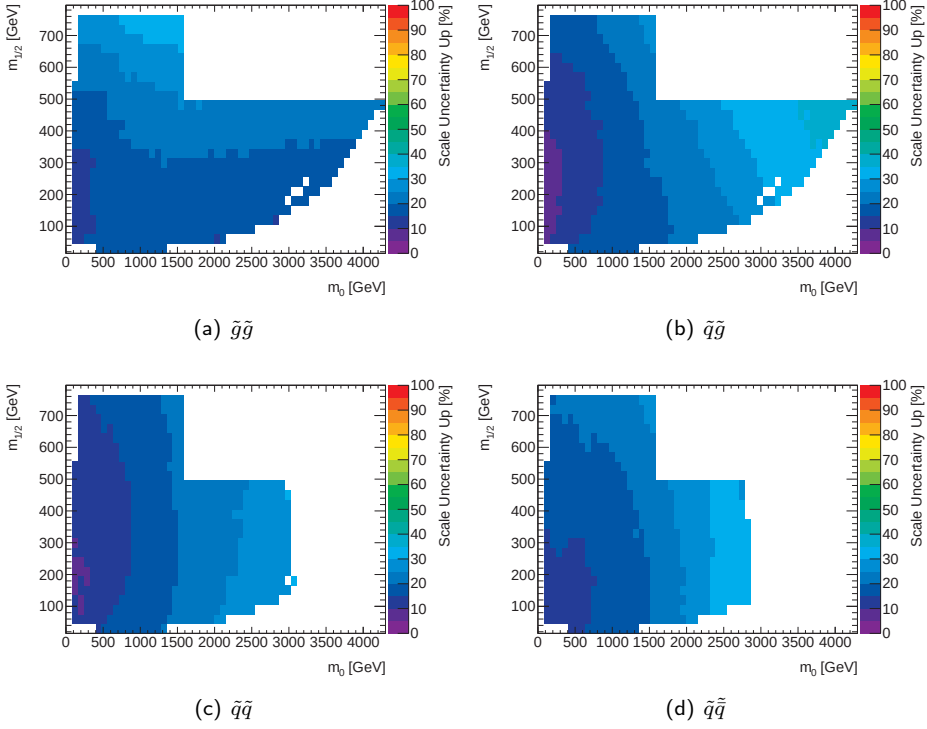


Figure 3.5: The relative upward scale uncertainty $\Delta_{\mu, CTEQ}^+$ on next-to-leading order cross sections calculated with PROSPINO using CTEQ6.6 PDFs in %. Shown are various production processes in the CMSSM plane at $\sqrt{s} = 7$ TeV: (a) gluino-pair production, (b) associated squark-gluino production, (c) squark-pair production and (d) squark-antisquark production.

we take here the most common definition by varying it between $\frac{1}{2}Q$ and $2Q$ [164]. We only vary μ while leaving all else constant, resulting in an two variations of the cross section, $\sigma_{\mu=Q/2}$ and $\sigma_{\mu=2Q}$. The up- and downward uncertainties on the cross section $\Delta_{\mu}^{\pm}$ are then given by

$$\Delta_{\mu}^+ = \frac{\max(\sigma_{\mu=Q/2} - \sigma, \sigma_{\mu=2Q} - \sigma)}{\sigma} \quad (3.8)$$

$$\Delta_{\mu}^- = \frac{\min(\sigma - \sigma_{\mu=Q/2}, \sigma - \sigma_{\mu=2Q})}{\sigma}. \quad (3.9)$$

Figures 3.5 and 3.6 show the up and down scale uncertainty of our four processes on NLO cross sections, again in the CMSSM plane. These uncertainties are mostly around $\sim 15\%$, yet for high m_0 they range up to 35%.

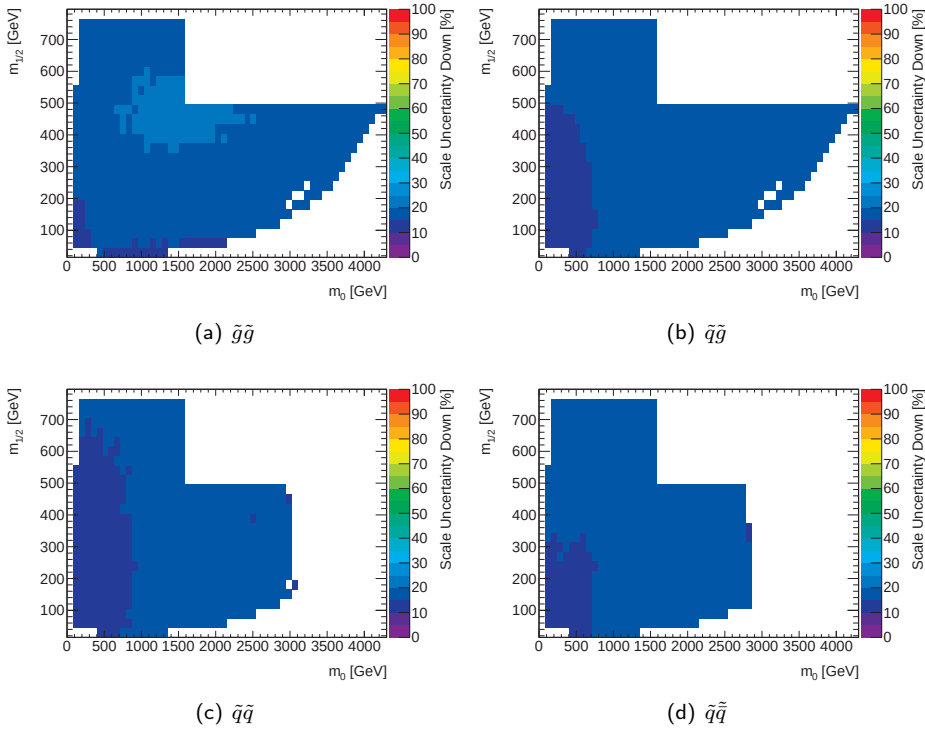


Figure 3.6: The relative downward scale uncertainty $\text{abs}(\Delta_{\mu, \text{CTEQ}}^-)$ on next-to-leading order cross sections calculated with PROSPINO using CTEQ6.6 PDFs in %. Shown are various production processes in the CMSSM plane at $\sqrt{s} = 7$ TeV: (a) gluino-pair production, (b) associated squark-gluino production, (c) squark-pair production and (d) squark-antisquark production.

Adding higher order terms to the calculation should reduce this scale dependence. This is seen in figure 3.7, where the dependence of the LO, NLO and NLO+NLL cross section as a function of μ is shown for gluino-pair and squark-pair production, for $m_{\tilde{g}} = m_{\tilde{q}} = 700$ GeV. The reduction in scale dependence between LO and NLO is huge. For gluino-pair production, the step to NLO+NLL calculations improves the dependence even further, while for squark-pair production the effect is much smaller. This difference is caused by the fact that the soft-gluon radiation has a larger effect for the gluinos than for squarks.

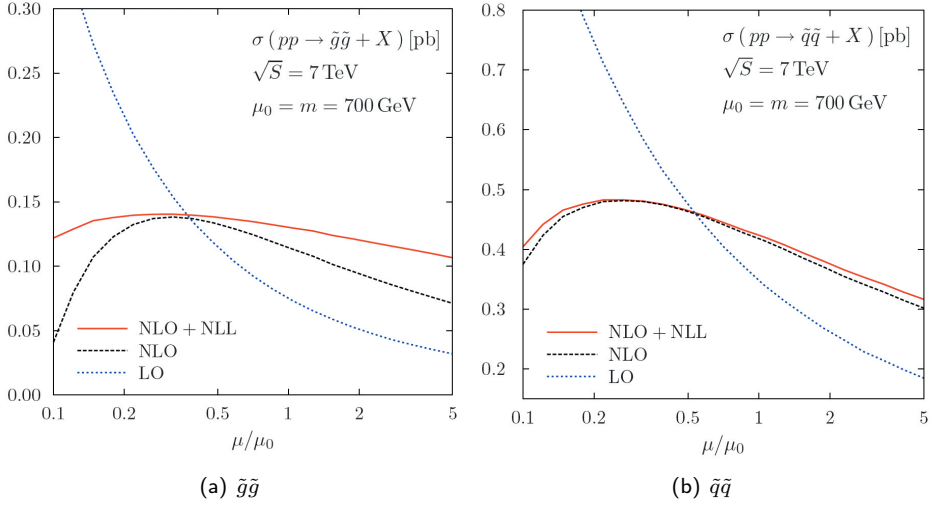


Figure 3.7: The dependence of the production cross section on the common factorisation and renormalisation scale μ , normalised to the nominal scale μ_0 . μ_0 is chosen as the mass of the gluino (left) or squark (right). Shown are leading order (LO), next-to-leading order (NLO) and next-to-leading order with soft gluon resummation (NLO+NLL) for (a) gluino-pair production and (b) squark-pair production at the LHC with $\sqrt{s} = 7$ TeV. Here the masses of the squarks and gluinos have been set equal. The figures were taken from [157].

3.4 Combining the contributions: two methods

Using the NLO or NLO+NLL calculations done by PROSPINO or NLL-FAST described in section 3.1 and the uncertainties on these described in section 3.3, there are several ways to combine these into a central value cross section and a total theoretical uncertainties thereon. Two methods are shortly described which were used in ATLAS analyses during the writing of this thesis.

3.4.1 Method 1: Straightforward usage

The easiest and most straightforward method is to use the central value of the cross section as returned by the NLO or NLO+NLL calculations. This is the most obvious method when considering only one PDF set, e.g. CTEQ. This was the case for the ATLAS analyses on the first part of the 7 TeV dataset (on those analyses approved before January 2012), see for instance [165–167]. Here the choice for CTEQ PDFs is given by the fact that these are the PDFs used in the simulation of the SUSY events.

The central value cross section is defined as the nominal NLO cross section cal-

culated by PROSPINO³. The relative uncertainties from the PDFs Δ_{PDF} , strong coupling uncertainty Δ_{α_S} and scale variation Δ_μ can then be added quadratically to obtain an asymmetric total uncertainty on the cross section per production process:

$$\sigma = \sigma_{NLO} \quad (3.10)$$

$$\Delta\sigma^+ = \sqrt{(\Delta_{PDF})^2 + (\Delta_{\alpha_S}^+)^2 + (\Delta_\mu^+)^2} \quad (3.11)$$

$$\Delta\sigma^- = \sqrt{(\Delta_{PDF})^2 + (\Delta_{\alpha_S}^-)^2 + (\Delta_\mu^-)^2}. \quad (3.12)$$

This central value cross section is nothing more than the cross section discussed in section 3.2, which results were shown in figure 3.1. Using the above equation and the uncertainties on the cross section from PDFs, α_S and the scale μ shown in figures 3.3, 3.4, 3.5 and 3.6 we can calculate the asymmetric total uncertainties on the cross sections. The upward uncertainty is shown in figure 3.8, while the downward uncertainty is very similar due to the fact that the total uncertainty is driven by the symmetric PDF uncertainty Δ_{PDF} .

The uncertainties have the same features as the cross sections. As the cross section of $\tilde{g}\tilde{g}$ production depends on the gluino mass, the uncertainty will also only depend on this mass, and similar for the other processes.

3.4.2 Method 2: Final official ATLAS and CMS agreement

One of the main issues with the previously defined method of obtaining SUSY particle production cross sections and their uncertainties is that just the PDFs from the CTEQ collaboration are used. Yet the various collaborations obtain these PDFs by fitting measurements, which is not done in a uniform way, leading to differences in the PDFs between the different collaborations. However, there is not one collaboration whose PDFs are preferred a priori: the recommendation from PDF4LHC [168] is to use all PDFs derived from the most global fits, using results from HERA, Tevatron and fixed-target experiments. The three collaborations which conform to this recommendation are CTEQ, MSTW and NNPDF with their CTEQ6.6, MSTW2008 and NNPDF2.0 PDFs. The cross section and uncertainties should then be calculated by combining results obtained from PDFs of each of these collaborations. Within ATLAS and CMS only the first two are used because of technical incompatibilities between PROSPINO and the NNPDF PDFs. It should be noted that both the CTEQ6.6 and MSTW2008 fits miss information from the most recent very accurate combined HERA datasets. These are included in the newest CTEQ fit (CT10) and will be included in the new MSTW fits, however since these are not used for ATLAS SUSY simulation, we use the older CTEQ6.6 set here.

When using multiple PDF sets, the choice of the central value cross section is no longer trivial, nor are its uncertainties. In the following the method for obtaining these is given, independent of the order of the calculation. This ‘envelope’ method was used in ATLAS results from January 2012 onwards; see for instance [169–171].

³In the mentioned analyses the NLO+NLL calculations were not yet available.

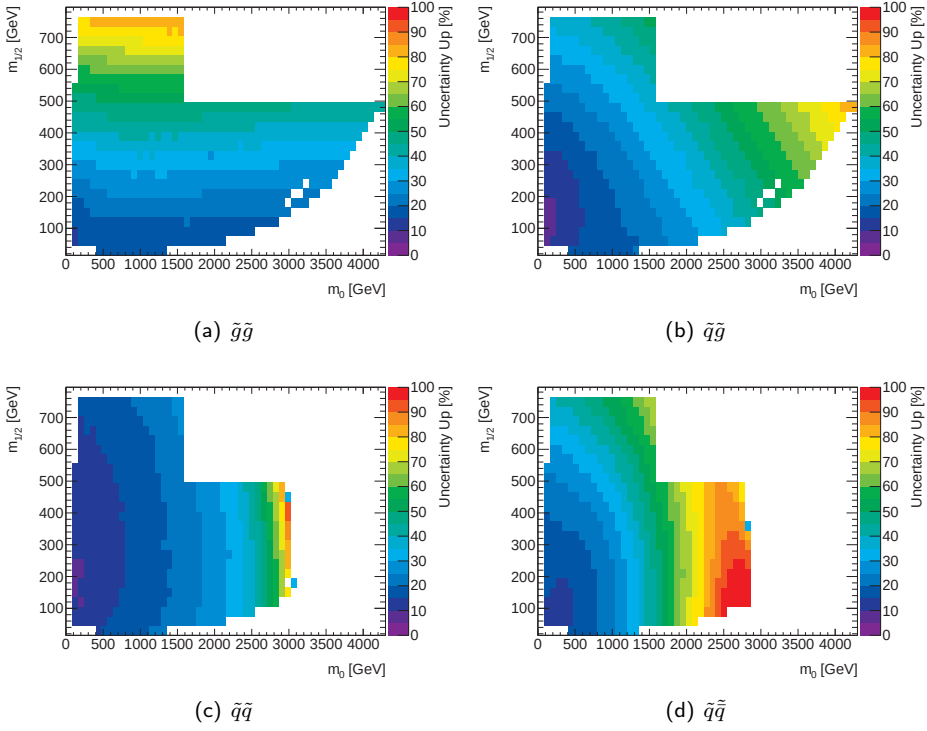


Figure 3.8: The total upward uncertainty Δ^+ on the NLO cross sections in %, where the PDF, scale and strong coupling uncertainties are combined. These are calculated with PROSPINO using CTEQ6.6 PDFs (the ‘straightforward’ method) for various production processes in the CMSSM plane at $\sqrt{s} = 7$ TeV: (a) gluino-pair production, (b) associated squark-gluino production, (c) squark-pair production and (d) squark-antisquark production.

Since none of the PDF sets is preferred a priori, the uncertainty on the cross section is taken as the maximum of the CTEQ and MSTW uncertainty bands, while the cross section is defined as the central value of this uncertainty band. More specifically: using the method described in sections 3.3.1 and 3.3.2 one can define the upward and downward one sigma variations due to PDF uncertainties (and α_S uncertainty for CTEQ PDFs) on the cross section for both PDF sets. Let $\delta_{PDF,CTEQ}^+$ ($\delta_{PDF,CTEQ}^-$) and $\delta_{PDF,MSTW}^+$ ($\delta_{PDF,MSTW}^-$) be the upward (downward) absolute variations of the cross section (in pb) using CTEQ and MSTW PDFs respectively, while the variations due to the strong coupling constant are called $\delta_{\alpha_S}^+$ and $\delta_{\alpha_S}^-$. Likewise the upward and downward variations due to factorisation and renormalisation scales (see section 3.3.3) are $\delta_{\mu,CTEQ}^+$ ($\delta_{\mu,CTEQ}^-$) and $\delta_{\mu,MSTW}^+$ ($\delta_{\mu,MSTW}^-$). These values define

a one sigma band for each uncertainty. By quadratically adding the three upward variations for CTEQ PDFs an upward CTEQ variation δ_{CTEQ}^+ is defined, etc.:

$$\delta_{\text{CTEQ}}^{\pm} = \sqrt{(\delta_{PDF, \text{CTEQ}}^{\pm})^2 + (\delta_{\mu, \text{CTEQ}}^{\pm})^2 + (\delta_{\alpha_s}^{\pm})^2}, \quad (3.13)$$

$$\delta_{\text{MSTW}}^{\pm} = \sqrt{(\delta_{PDF, \text{MSTW}}^{\pm})^2 + (\delta_{\mu, \text{MSTW}}^{\pm})^2}, \quad (3.14)$$

$$(3.15)$$

The maximum of the upward variations and minimum of the downward variations define a band around the nominal cross sections, σ_{CTEQ} and σ_{MSTW} , with upper and lower values defined by:

$$\sigma_+ = \max(\sigma_{\text{CTEQ}} + \delta_{\text{CTEQ}}^+, \sigma_{\text{MSTW}} + \delta_{\text{MSTW}}^+) \quad (3.16)$$

$$\sigma_- = \min(\sigma_{\text{CTEQ}} - \delta_{\text{CTEQ}}^-, \sigma_{\text{MSTW}} - \delta_{\text{MSTW}}^-). \quad (3.17)$$

The central value cross section is straightforwardly defined from the above as the centre of this envelope, with the absolute uncertainty being half of the width of the band:

$$\sigma = \frac{\sigma_+ + \sigma_-}{2} \quad (3.18)$$

$$\delta\sigma = \frac{\sigma_+ - \sigma_-}{2}. \quad (3.19)$$

The resulting uncertainty band is symmetric, as opposed to the previous method where an asymmetric uncertainty could be obtained. The relative uncertainty $\Delta\sigma$ is obtained by just dividing by the cross section, $\Delta\sigma = \delta\sigma/\sigma$. Note that using this method it is possible, with large differences between $\delta_{\text{CTEQ}}^{\pm}$ and $\delta_{\text{MSTW}}^{\pm}$, to have a final obtained cross section which is larger than $\max(\sigma_{\text{CTEQ}}, \sigma_{\text{MSTW}})$, or smaller than $\min(\sigma_{\text{CTEQ}}, \sigma_{\text{MSTW}})$. However, the method ensures that both nominal values fall safely in the systematic uncertainty envelope.

Results of this method are shown in figures 3.9 and 3.10 for the same 7 TeV CMSSM scenario as the previous method, which were shown in figures 3.1 and 3.8. It is apparent that the cross sections are similar, yet a little bit higher in some regions, to the previous method. The uncertainty on the cross section is similar as well: while the decreased scale uncertainty on the NLO+NLL calculations lowers the total uncertainty, the difference between σ_{CTEQ} and σ_{MSTW} cancels this effect for most processes.

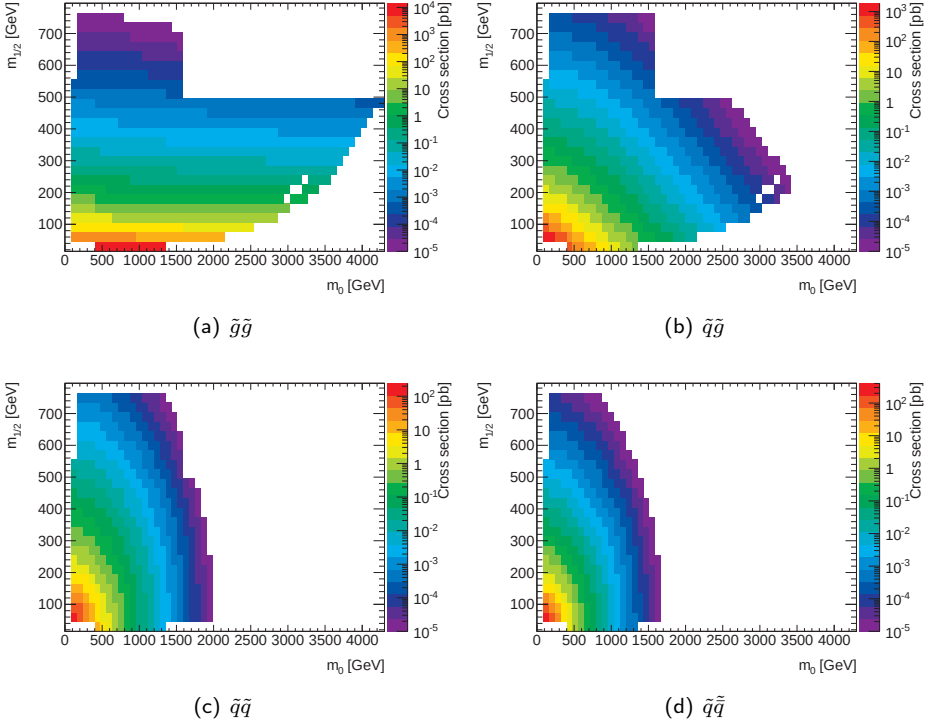


Figure 3.9: *The NLO+NLL cross sections derived using the ‘envelope’ method for various production processes in the CMSSM plane at $\sqrt{s} = 7$ TeV: (a) gluino-pair production, (b) associated squark-gluino production, (c) squark-pair production and (d) squark-antisquark production.*

3.4.3 Comparison of methods

To compare the methods in more details, it is easiest to concentrate at two points in the CMSSM plane and view the various contributions separately. In tables 3.1 and 3.2 the differences between the two methods are shown for two CMSSM models on different parts of the plane, for the production of the four possible coloured SUSY pairs ($\tilde{g}\tilde{g}$, $\tilde{q}\tilde{g}$, $\tilde{q}\tilde{q}$ and $\tilde{q}\tilde{\bar{q}}$) at the LHC running at $\sqrt{s} = 8$ TeV. The point with $m_0 = 600$ GeV and $m_{1/2} = 700$ GeV has squark and gluino mass very close together (~ 1550 GeV), while the second point with $m_0 = 2600$ and $m_{1/2} = 350$ has high squark mass (~ 2.2 TeV) and relatively light gluino mass (~ 900 GeV). These points lie on the expected exclusion limit in the CMSSM plane of the 5.8 fb^{-1} 8 TeV full hadronic analysis [172]. The high squark mass in the second point is outside of the interpolation reach of NLL-FAST for the squark production, where no soft gluon corrections can be applied, and thus leads to NLO values for the squark-pair and squark-antisquark production.

These tables show in the last row the cross section in picobarn with its uncertainty,

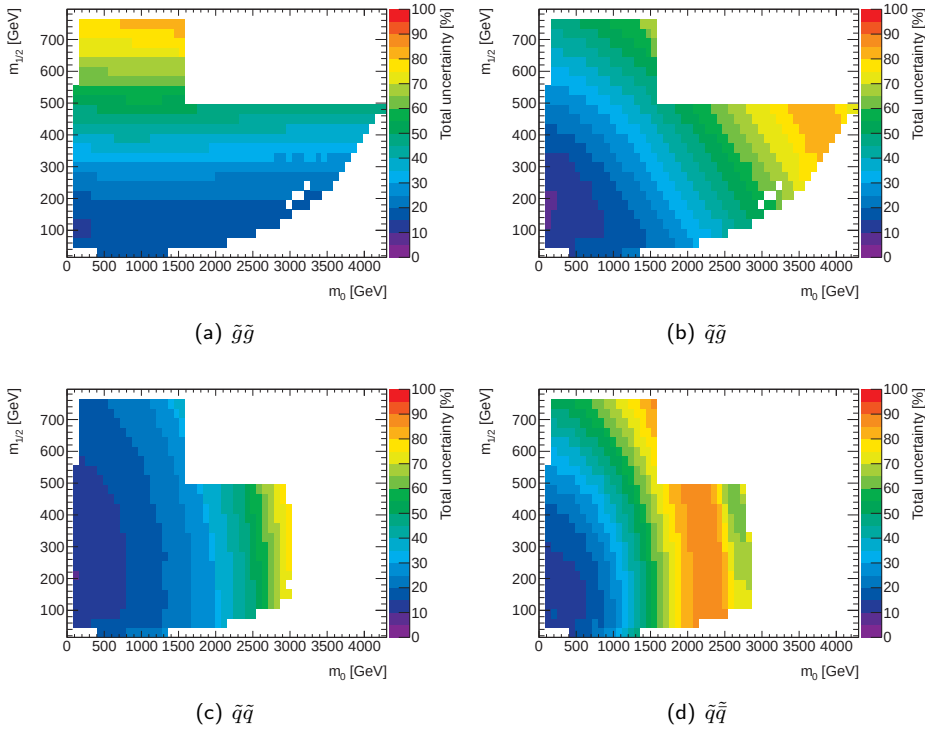


Figure 3.10: The relative uncertainties in % on the NLO+NLL cross sections derived using the ‘envelope’ method for various production processes in the CMSSM plane at $\sqrt{s} = 7$ TeV: (a) gluino-pair production, (b) associated squark-gluino production, (c) squark-pair production and (d) squark-antisquark production.

for various contributions to this cross section and uncertainty. The cross section is seen to be higher for all processes in table 3.2, which is due to the addition of the NLL order diagrams in the calculations. This can be seen from the first row, which shows the nominal CTEQ cross section. Contrary to the expectations, the uncertainties have increased. As already pointed out, a decrease in the scale uncertainties due to addition of the soft gluon resummation does not mitigate the effect of the difference between $\sigma_{nom,CTEQ}$ and $\sigma_{nom,MSTW}$, which is seen to be large. For sub-processes involving a gluino, $\Delta_{PDF,CTEQ}$ dominate the total uncertainty, while for the squark sub-processes the scale uncertainties contribute significantly. The contribution of δ_{α_S} are small, justifying not including MSTW α_S uncertainties.

$m_0 = 600 \text{ GeV}, m_{1/2} = 700 \text{ GeV}$				
	$\tilde{g}\tilde{g}$	$\tilde{q}\tilde{g}$	$\tilde{q}\tilde{q}$	$\tilde{q}\tilde{\bar{q}}$
$\sigma_{NLO, \text{CTEQ}} [\text{pb}]$	8.23×10^{-5}	0.00107	0.00197	0.000221
$\Delta_{PDF, \text{CTEQ}} [\%]$	55.1	24.8	5.31	25.9
$\Delta_{\mu, \text{CTEQ}}^+ [\%]$	22.7	15	13.7	19
$\Delta_{\mu, \text{CTEQ}}^- [\%]$	-17.5	-15.3	-14.7	-16.3
$\Delta_{\alpha_S, \text{CTEQ}}^+ [\%]$	14.6	5.68	2.16	6.88
$\Delta_{\alpha_S, \text{CTEQ}}^- [\%]$	-9.38	-3.41	-2.16	-3.85
$\Delta^+ \sigma [\%]$	61.4	29.5	14.9	32.8
$\Delta^- \sigma [\%]$	-58.6	-29.4	-15.8	-30.8
$\sigma [\text{pb}]$	$8.23^{(+5.05)}_{(-4.82)} \times 10^{-5}$	$1.07^{(+0.32)}_{(-0.31)} \times 10^{-3}$	$1.97^{(+0.29)}_{(-0.31)} \times 10^{-3}$	$2.21^{(+0.73)}_{(-0.68)} \times 10^{-4}$
$m_0 = 2600 \text{ GeV}, m_{1/2} = 350 \text{ GeV}$				
	$\tilde{g}\tilde{g}$	$\tilde{q}\tilde{g}$	$\tilde{q}\tilde{q}$	$\tilde{q}\tilde{\bar{q}}$
$\sigma_{NLO, \text{CTEQ}} [\text{pb}]$	0.0381	0.000396	2.2×10^{-7}	2.17×10^{-8}
$\Delta_{PDF, \text{CTEQ}} [\%]$	22.4	29.9	20.8	73.7
$\Delta_{\mu, \text{CTEQ}}^+ [\%]$	18.6	24.2	22.7	25.8
$\Delta_{\mu, \text{CTEQ}}^- [\%]$	-17.8	-16.4	-15	-15.2
$\Delta_{\alpha_S, \text{CTEQ}}^+ [\%]$	5.9	6.29	2.76	16.8
$\Delta_{\alpha_S, \text{CTEQ}}^- [\%]$	-3.83	-4.3	-4.42	-11.8
$\Delta^+ \sigma [\%]$	29.8	39	30.9	79.8
$\Delta^- \sigma [\%]$	-28.9	-34.4	-26	-76.1
$\sigma [\text{pb}]$	$0.0381^{(+0.0113)}_{(-0.011)}$	$3.96^{(+1.55)}_{(-1.36)} \times 10^{-4}$	$2.2^{(+0.68)}_{(-0.57)} \times 10^{-7}$	$2.17^{(+1.73)}_{(-1.65)} \times 10^{-8}$

Table 3.1: The NLO cross section at $\sqrt{s} = 8 \text{ TeV}$ in pb and its uncertainty in brackets, together with all components of the uncertainty for the straightforward method, where only CTEQ uncertainties are used. Given are the values for the most important processes of two CMSSM models, with in the upper table a point which has $m_0 = 600 \text{ GeV}$, $m_{1/2} = 700 \text{ GeV}$, $\tan\beta = 10$, $A_0 = 0$ and positive μ , while the lower has $m_0 = 2600 \text{ GeV}$, $m_{1/2} = 350 \text{ GeV}$.

$m_0 = 600 \text{ GeV}, m_{1/2} = 700 \text{ GeV}$				
	$\tilde{g}\tilde{g}$	$\tilde{q}\tilde{g}$	$\tilde{q}\tilde{q}$	$\tilde{q}\tilde{\bar{q}}$
$\sigma_{nom, \text{CTEQ}} [\text{pb}]$	1.05×10^{-4}	0.00123	0.00206	2.41×10^{-4}
$\sigma_{nom, \text{MSTW}} [\text{pb}]$	5.02×10^{-5}	0.000861	0.0019	1.70×10^{-4}
$\Delta_{PDF, \text{CTEQ}}^+ [\%]$	69	29.6	5.8	32.1
$\Delta_{PDF, \text{MSTW}}^+ [\%]$	26.1	11.4	3.9	13
$\Delta_{PDF, \text{CTEQ}}^- [\%]$	-41.3	-20.7	-5.6	-20.6
$\Delta_{PDF, \text{MSTW}}^- [\%]$	-23.6	-10.8	-2.9	-12.6
$\Delta_{\mu, \text{CTEQ}}^+ [\%]$	5.7	6.02	10.7	13.8
$\Delta_{\mu, \text{MSTW}}^+ [\%]$	2.93	5.68	10.9	14.8
$\Delta_{\mu, \text{CTEQ}}^- [\%]$	-7.87	-9.02	-11.7	-12.7
$\Delta_{\mu, \text{MSTW}}^- [\%]$	-5.2	-8.55	-11.9	-13.3
$\Delta_{\alpha_S, \text{CTEQ}}^+ [\%]$	16.6	7.2	0.2	9.1
$\Delta_{\alpha_S, \text{CTEQ}}^- [\%]$	-10.8	-5.3	-0.5	-5.7
$\Delta^\pm \sigma [\%]$	65.5	37.1	16.2	40.7
$\sigma [\text{pb}]$	$1.09(\pm 0.71) \times 10^{-4}$	$1.18(\pm 0.44) \times 10^{-3}$	$1.99(\pm 0.32) \times 10^{-3}$	$2.33(\pm 0.95) \times 10^{-4}$
$m_0 = 2600 \text{ GeV}, m_{1/2} = 350 \text{ GeV}$				
	$\tilde{g}\tilde{g}$	$\tilde{q}\tilde{g}$	$\tilde{q}\tilde{q}$	$\tilde{q}\tilde{\bar{q}}$
$\sigma_{nom, \text{CTEQ}} [\text{pb}]$	0.0422	4.50×10^{-4}	2.2×10^{-7}	2.17×10^{-8}
$\sigma_{nom, \text{MSTW}} [\text{pb}]$	0.0371	2.74×10^{-4}	1.43×10^{-7}	4.03×10^{-9}
$\Delta_{PDF, \text{CTEQ}}^+ [\%]$	26.5	36	25.3	97.9
$\Delta_{PDF, \text{MSTW}}^+ [\%]$	11	13.6	4.42	29.8
$\Delta_{PDF, \text{CTEQ}}^- [\%]$	-18.3	-24.5	-17.6	-50.1
$\Delta_{PDF, \text{MSTW}}^- [\%]$	-10.9	-12.4	-6.71	-28.9
$\Delta_{\mu, \text{CTEQ}}^+ [\%]$	10.4	14.7	22.7	25.8
$\Delta_{\mu, \text{MSTW}}^+ [\%]$	10.8	13.7	23.1	18.9
$\Delta_{\mu, \text{CTEQ}}^- [\%]$	-10	-13.4	-15	-15.2
$\Delta_{\mu, \text{MSTW}}^- [\%]$	-10.2	-12.4	-15.4	-15.1
$\Delta_{\alpha_S, \text{CTEQ}}^+ [\%]$	8	8.3	2.76	16.8
$\Delta_{\alpha_S, \text{CTEQ}}^- [\%]$	-5.5	-6.1	-4.42	-11.8
$\Delta^\pm \sigma [\%]$	27.1	47.3	42.5	88.4
$\sigma [\text{pb}]$	$0.043(\pm 0.011)$	$4.27(\pm 2.02) \times 10^{-4}$	$2.07(\pm 0.88) \times 10^{-7}$	$2.33(\pm 2.06) \times 10^{-8}$

Table 3.2: The NLO+NLL cross section at $\sqrt{s} = 8 \text{ TeV}$ in pb with its uncertainty in brackets, together with all components of the uncertainty for method 2. Given are the values for the most important processes of two CMSSM models, with in the upper table a point which has $m_0 = 600 \text{ GeV}$, $m_{1/2} = 700 \text{ GeV}$, $\tan\beta = 10$, $A_0 = 0$ and positive μ , while the lower has $m_0 = 2600 \text{ GeV}$, $m_{1/2} = 350 \text{ GeV}$. Note that for the second point the nominal cross sections σ_{nom} of the two gluino processes, $\tilde{g}\tilde{g}$ and $\tilde{q}\tilde{g}$, are NLO+NLL values, while the two squark-production channels are only at NLO.

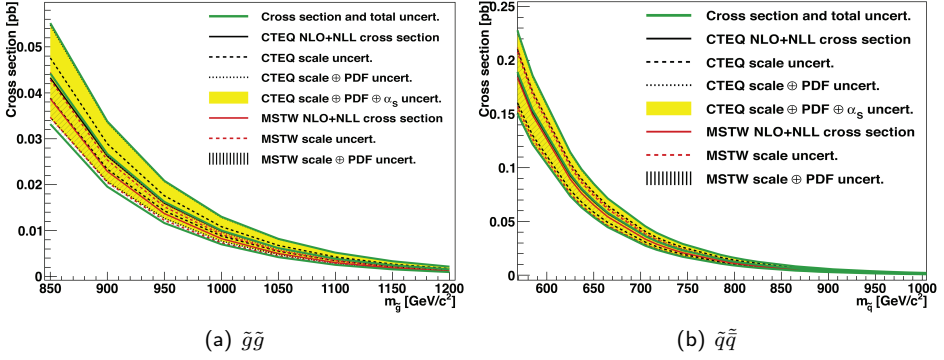


Figure 3.11: The NLO+NLL cross section for 2 simplified models: (a) gluino-gluino production, where the squarks are decoupled, as a function of gluino mass, and (b) squark-antisquark production as a function of squark mass, where the gluinos and third generation squarks are decoupled. The cross section is shown by the middle green line, with the total uncertainty given by the outer two green lines. The solid lines show the CTEQ6.6 (black) and MSTW2008 (red) NLO+NLL cross sections. The dashed lines show the scale uncertainty, while the dotted black line indicates the uncertainty due to the scale and PDF uncertainty using CTEQ. The yellow band and black dashed regions show the total uncertainty on the CTEQ and MSTW cross section, respectively [173].

3.4.4 Method for specific models

This chapter has mostly discussed the CMSSM, yet as discussed in section 1.3.4, many analyses interpret their results using simplified models, where only a certain number of particles can be produced at the LHC. For these simplified models the method of calculating the cross sections is not different from the general SUSY model. However, by decoupling certain particles from the rest of the phenomenology at the LHC energy scales, some processes will no longer be allowed, and others will be slightly altered. Below we will first go through some of these simplified models, and finally show results for the electroweak sector.

Gluino-pair production

To obtain cross sections for pure gluino-pair production⁴ the squarks are completely decoupled using NLL-FAST, i.e. diagrams including squarks do not contribute to $\sigma_{\tilde{g}\tilde{g}}$. This cross section is shown in figure 3.11(a), where the green solid lines show the central value and its uncertainty band. The various components are also shown:

⁴In this thesis the main focus is on models with the direct decay of the sparticles, i.e. of the gluino decaying into two quarks and a lightest neutralino, $\tilde{g} \rightarrow q + q + \tilde{\chi}_1^0$. Yet the cross section does not depend on the decay modes: as long as the squark is decoupled the gluino-pair production cross section will be the same.

in black (red) the NLO+NLL nominal cross section and its uncertainties from the factorisation and renormalisation scale uncertainty, the PDF uncertainty and the strong coupling uncertainty for the CTEQ (MSTW) PDF sets. The yellow band corresponds to the total uncertainty using CTEQ PDFs, while the dashed region corresponds to the total MSTW uncertainty. The gluino mass range is chosen to show the sensitivity of ATLAS and CMS with their $\sqrt{s} = 7$ TeV full dataset analyses.

From this figure it can be seen that while the cross sections derived with MSTW PDFs are lower than those predicted by CTEQ PDFs, the MSTW uncertainty band falls nearly completely inside the CTEQ uncertainty band. In the mass range shown the uncertainty using CTEQ PDFs is approximately twice that of the uncertainty using MSTW PDFs. This means that the MSTW contribution to the green ‘envelope’, and thus to the central value of the cross section, is nearly negligible.

Squark-antisquark production

Many simplified models which focus on just squark production, decaying into squarks and neutralinos, have decoupled gluinos and third generation squarks. As the t-channel squark-pair production is not possible without gluinos (see section 1.3.4), only squark-antisquark production is allowed.

Another important issue are the third generation squarks. In many simplified models the first two generation squarks have a degenerate mass, while the third generation is decoupled (see for instance the simplified models introduced in section 1.3.4). Yet the calculations from both PROSPINO and NLL-FAST take all five light-flavour squarks, including the sbottoms, to be degenerate in mass, resulting in a cross section which is too high. As all degenerate squarks contribute equally, the cross section is rescaled by a factor of 4/5 to remove the sbottom contribution.

Other simplified models exist (e.g. those focussing on decays via gauginos) which don’t only decouple the third generation but also the right-handed squarks. With 10 final states contributing to the calculated cross section ($\tilde{q}\tilde{q}$, with $\tilde{q} = (\tilde{q}_L, \tilde{q}_R$ and q one of the five light-flavour squarks), and only 4 final states requested, the cross section needs to be rescaled by a factor 4/10.

In figure 3.11(b) the results are shown in the same fashion as the gluino-pair production results. The CTEQ and MSTW central value cross sections are very close together, yet the MSTW uncertainty band is somewhat smaller, making the result mostly depending on the CTEQ numbers.

Third generation squark production

To study the impact of ATLAS analyses on the third generation squarks, some simplified models have been constructed where only two lightest top squarks, or only two lightest bottom squarks are produced [174, 175]. The cross sections for these are equal at leading order – at higher orders small differences are introduced from e.g. the stop/sbottom mixing angle. However, these effects are negligible [152], allowing us to calculate the stop and sbottom cross section identically (for the same mass and mixing parameters).

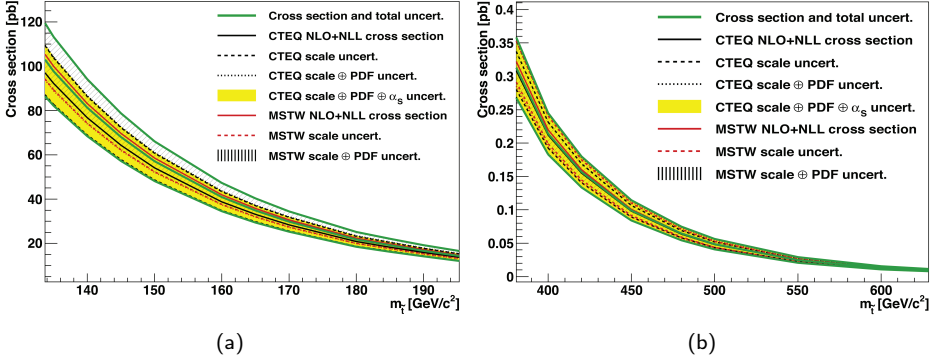


Figure 3.12: The NLO+NLL cross section for direct-stop production, where the gluinos and first two generation squarks are decoupled, as a function of the stop mass: (a) shows low stop masses, (b) high stop masses. The cross section is shown by the middle green line, with the total uncertainty given by the outer two green lines. The black solid line shows the CTEQ6.6 NLO+NLL cross section, while the black dashed line shows the scale uncertainty, the dotted line the uncertainty due to the scale and PDF uncertainty and the yellow band shows the total uncertainty on the CTEQ cross section. In red the MSTW2008 NLO+NLL cross section is shown, with the dashed red line its scale uncertainty, and the dashed region its total scale + PDF uncertainty [173].

Figure 3.12 shows the results for this scenario. In the left figure, it is seen that for low stop masses the prediction for the cross section using the MSTW PDF set yield a higher central value than using the CTEQ set, with as usual smaller uncertainties. This leads to a significantly higher cross section and uncertainty than using the CTEQ set alone. When looking at higher masses in the right figure, the results for the CTEQ and MSTW PDF set are very similar, with slightly larger uncertainties for the former set. Thus the prediction of the cross section is largely due to the CTEQ PDF set.

Production of electroweak SUSY particles

In this chapter the production of electroweak SUSY particles (gauginos and sleptons) has not been studied yet extensively, as these are not produced significantly at the LHC. Yet with the rapidly increasing luminosity at the LHC, the direct production of electroweak SUSY particles is starting to be important. Figure 3.13 shows the *inclusive* 7 TeV NLO cross section using the ‘envelope’ method for gaugino-pair production ($\tilde{\chi}\tilde{\chi}$), the associated production of a gaugino with a gluino or a squark ($\tilde{\chi}\tilde{g}$ or $\tilde{\chi}\tilde{q}$), and finally slepton pairs ($\tilde{l}\tilde{l}$), for a CMSSM scenario. Here the contributions of all possible combinations of gauginos or sleptons are summed. The cross sections are smaller than

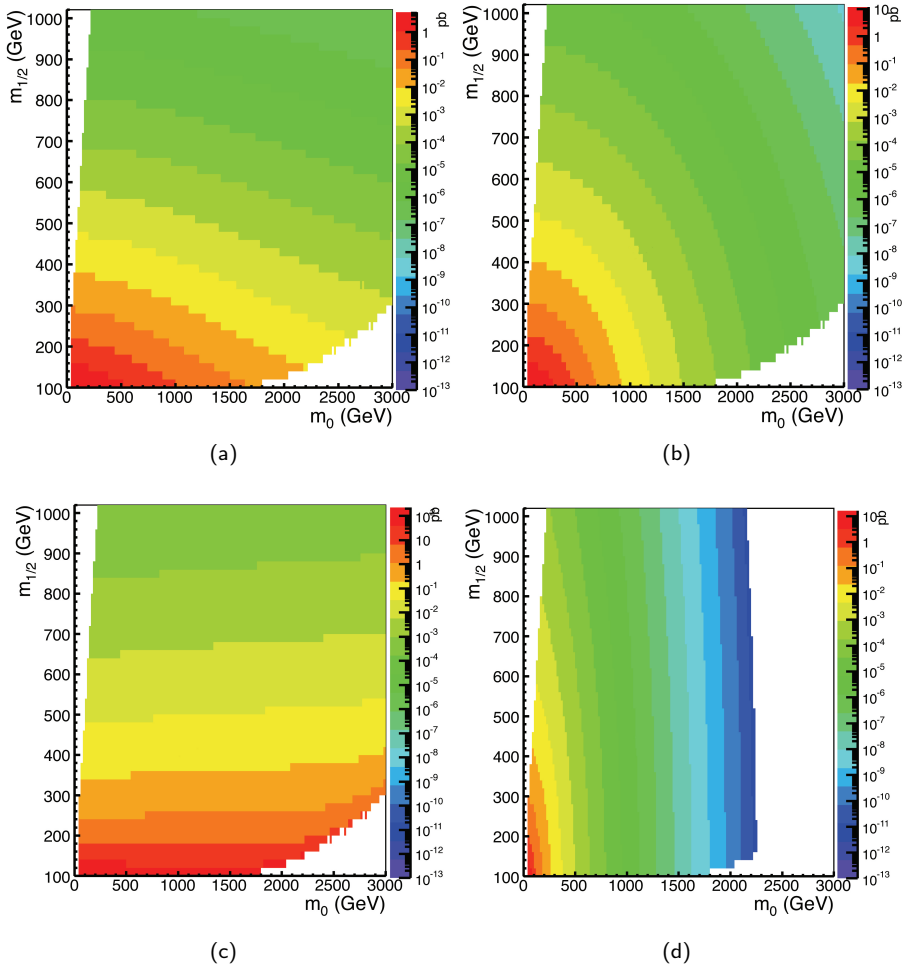


Figure 3.13: The NLO cross section for various electroweak processes in the CMSSM plane at $\sqrt{s} = 7$ TeV. Top left (right) shows the cross section for inclusive gaugino-gluino (squark) production. Bottom left shows inclusive gaugino-pair production, while bottom right shows inclusive slepton-pair production [173].

for the coloured sector, except for gaugino-pair production, where the cross section for low $m_{1/2}$ is comparable to coloured sparticle production. The uncertainty on $\sigma_{\tilde{\chi}\tilde{\chi}}$ is around 10%, while on $\sigma_{\tilde{\chi}\tilde{q}}$ and $\sigma_{\tilde{\chi}\tilde{g}}$ the uncertainties are around 35%, mainly because of PDF uncertainties due to the coloured objects. Uncertainties on pure slepton production are negligible ($< 5\%$).

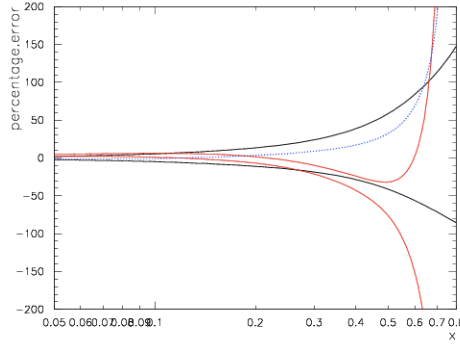


Figure 3.14: *Uncertainty on the gluon PDF as a function of the Björken scaling x , for CTEQ6.6 (in black), MSTW2008NLO (in blue) and NNPDF2.2NLO (in red). The uncertainty band of the CTEQ set covers all 3 sets up to $x < 0.3$. Above this, NNPDF can become smaller than CTEQ, making it important to include. This plot was made using [176].*

3.4.5 Discussion

The procedure discussed in section 3.4.2 still does not yet fully comply with the PDF4LHC recommendations, as the PDFs from the NNPDF collaboration has not yet been implemented due to technical issues. Yet this might have an impact on the SUSY cross section predictions, as seen in figure 3.14. This shows the dependence of the central value of a PDF and its uncertainty band, for the three PDF sets on the Björken scaling x [177]. The black CTEQ6.6 band is seen to cover the other two PDFs very well for low x , yet the red band of the NNPDF2.2NLO set is seen to be quite a bit smaller than the other two for high x . This means that as one probes higher energies (higher x means a more energetic initial parton) it is important to include the NNPDF2.2NLO set in the calculation of the production cross sections.

Event reconstruction and simulation in ATLAS

Particles produced in proton-proton collisions in the centre of ATLAS are detected by the sub-detectors described in chapter 2. To identify which particles traversed the detector, the information from the sub-detectors needs to be analysed and reconstructed. From all types of Standard Model particles produced in pp collisions, only neutrinos do not interact with the detector, while other leptons, hadrons and photons are all detected by ATLAS. These particles need to be reconstructed and identified. The reconstruction procedure of detected particles is described in sections 4.1-4.4, resulting in physics objects which can be used in the analyses described in the following chapters. Although neutrinos do not give a signal in the detector, knowing that the net momentum in the transverse plane should be zero due to momentum conservation, an momentum imbalance in this transverse plane can be detected when an energetic neutrino passes through the detector. This imbalance is called the *missing transverse energy*, or $\cancel{E}_T$, where the step from momentum to energy is taken by assuming negligible mass. Its reconstructing is described in section 4.5.

To compare observations with theoretical expectations, data simulation is another extremely important aspect of these analyses. Event generation using Monte Carlo techniques, together with the simulation of the response of the ATLAS detector, is discussed in section 4.7.

4.1 Track and vertex reconstruction

The reconstruction of physics objects starts in many cases with identifying tracks and vertices in the inner detector. Tracks are found using the space point measurements from the pixel, SCT and TRT sub-detectors. Using the NEWT tracking reconstruction algorithm [178], charged particles are traced from the inner detector to the outer detectors (inside-out tracking): first, hits in the silicon detectors are used. Two-dimensional space point measurements are obtained directly from the pixel detectors. The silicon strips of the SCT on the other hand can only provide information in one

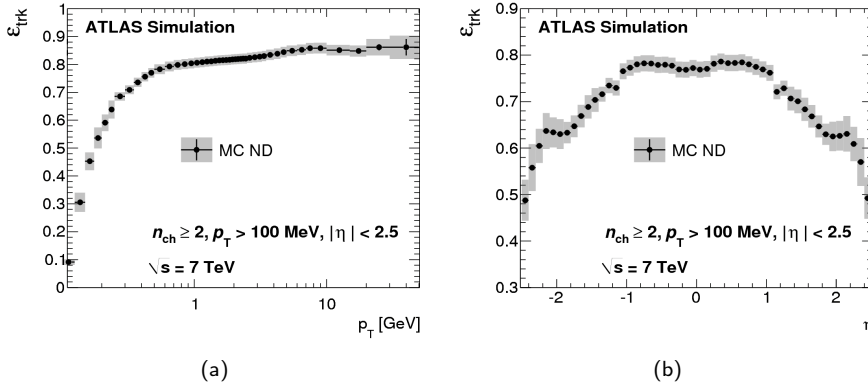


Figure 4.1: Track reconstruction efficiency as a function of (a) p_T and (b) η , measured in non-diffractive MC simulation for events with at least 2 charged particles in the event. Statistical uncertainties are represented by the black lines, while the green band includes the systematics. Taken from [180].

direction. As described in chapter 2, SCT modules consist of two sides placed on top of each other, with an angle of 40 mrad between them. A hit on both sides is therefore needed to obtain a two-dimensional space point measurement. Finally, the third space coordinate is taken as the position of the module.

Once three space points are found in separate layers of the pixel and SCT sub-detectors, they form a seed for the tracking algorithm. From the direction of the seed track, a ‘road’ is defined where the hits are expected. From within the road hits in the other layers of the silicon detector are gathered using a Kalman filter [179]. Ambiguities, from for instance fake tracks or overlapping segments, are solved by giving priority to measurements with more precision. Tracks are then extended into the TRT. On top of these ‘default’ tracks, a more ‘robust’ set of requirements can be applied to minimise the effect of fake tracks. This is achieved by requiring tracks to have at least 9 hits and no holes in the pixel detector. Here a ‘hole’ is space point where a hit was expected yet not detected.

For each reconstructed track the distance to the z -axis (d_0) and the z -coordinate at the point of closest approach to the z -axis (z_0) are determined from a global fit to all assigned hits. d_0 and z_0 called the impact parameters of the track.

The efficiency of the tracking algorithm is defined by the fraction of truth particles¹ with $p_T > 100$ MeV and $|\eta| < 2.5$ which can be matched to a reconstructed track. Figure 4.1 shows this efficiency for particles taken from Monte Carlo (MC, see section 4.7), as a function of p_T and η of the track. For high p_T tracks, the track reconstruction efficiency is close to 90%. The reconstruction algorithm is checked to be resilient against pile-up for tracks with $p_T > 10$ GeV [141].

Using the reconstructed tracks, the vertices from which these tracks originate can

¹Truth particles in simulated events are particles before detector simulation

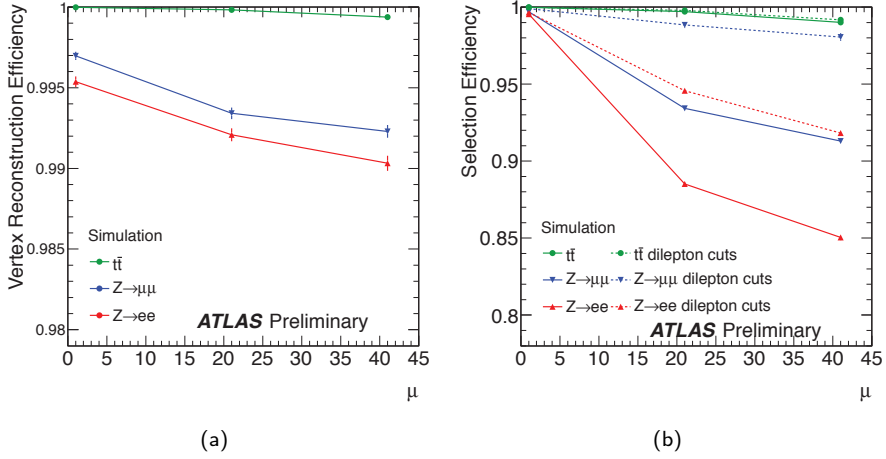


Figure 4.2: Efficiency for (a) vertex reconstruction and (b) reconstructing the hard scattering vertex and then select it as primary vertex, as a function of the average number of interactions per event, μ , measured in MC simulation of $t\bar{t}$, $Z \rightarrow e^+e^-$ and $Z \rightarrow \mu^+\mu^-$ events. Solid lines are efficiencies before any event selection, dashed lines are after requiring two leptons. Taken from [183, 184].

be reconstructed, leading to one or more pp interaction points (primary vertices), and secondary vertices belonging to particle decays. This vertex reconstruction is performed by first applying a vertex finding algorithm, associating reconstructed tracks to vertex candidates, after which a vertex fitting algorithm is performed reconstructing the position of the vertex precisely.

Reconstructed tracks are selected if they are coming from the interaction region, which is determined during the physics run [181]. From this selection, a vertex seed is found by searching for the most often occurring z -position of the tracks. The exact position of the vertex is determined using a χ^2 based vertex fitting algorithm, using the seed and the tracks around it, where the contribution of outlying tracks is decreased with respect to non-outliers [182]. Tracks which are incompatible with the vertex by at least 7σ seed new vertices, until no additional vertex can be found with at least 2 tracks. The primary vertex is finally taken to be the vertex with the highest $\sum p_T^2$ of tracks.

The vertex reconstruction efficiency in MC simulation of $t\bar{t}$, $Z \rightarrow e^+e^-$ and $Z \rightarrow \mu^+\mu^-$ events is shown in figure 4.2 (a), as a function of the average number of pp collisions per bunch crossing. The efficiencies are calculated before any data analysis event selection was applied. The vertex reconstruction efficiency is above 99% for all three processes, with a slight decrease in efficiency for higher μ , as expected. Figure 4.2 (b) shows the efficiency to reconstruct the hard scattering vertex, and subsequently select it as the primary vertex for the same processes, with and without

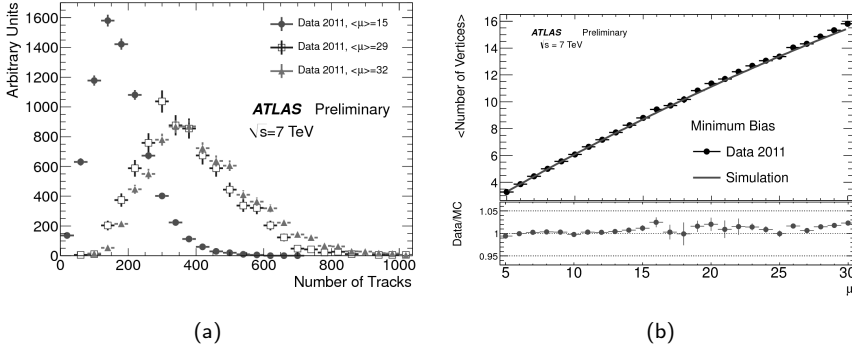


Figure 4.3: (a) Comparison of the number of reconstructed tracks per event between datasets with varying average number of interactions per event. (b) Average number of primary vertices, reconstructed with default tracks, as a function of the average number of interactions $\langle\mu\rangle$ per event, compared to events simulated with PYTHIA. Both plots are made from data taken in 2011. Taken from [141, 183].

selecting dilepton events. In $t\bar{t}$ events the selection efficiency is above 99%. After the event selection performed for the SUSY analysis described in chapter 5, events are required to have a well reconstructed vertex with at least 5 tracks associated to it. The efficiency of this requirement is 98.6% in the total dataset of 5.8 fb^{-1} at $\sqrt{s} = 8 \text{ TeV}$, while in $t\bar{t}$ events simulated with POWHEG+PYTHIA it is again above 99.9%. The efficiency of the cut is studied in SUSY signal events by checking several signal samples of the squark-pair and gluino-pair simplified models, and is observed to be 100% for each of the models.

The vertex position resolution, depending on the number of tracks associated to it, ranges from 0.6 mm for less than 5 tracks to 0.03 mm for > 80 tracks in the z direction, and between 0.35 mm and 0.02 mm in the x direction. This is calculated from a $\sqrt{s} = 8 \text{ TeV}$ dataset with a random trigger, in a dedicated fill with a very low average number of pp collisions per bunch crossing of only 0.01 [184].

The number of reconstructed tracks per event obviously depends heavily on pile-up. Figure 4.3 (a) shows a distribution of the number of reconstructed tracks per event in data recorded in 2011 at $\sqrt{s} = 7 \text{ TeV}$, with three different values of the mean number of collisions per bunch crossing $\langle\mu\rangle$. With increasing $\langle\mu\rangle$, the average number of reconstructed tracks significantly. The average number of reconstructed primary vertices N_{PV} is shown in figure 4.3 (b) as a function of $\langle\mu\rangle$, comparing 2011 data collected using a minimum bias trigger with a dataset simulated with PYTHIA.

4.2 Jets

Just like in many Standard Model processes, the production of squarks and gluinos and their subsequent decays leads to highly energetic quarks and gluons flying away from the interaction point. Both quarks and gluons carry colour charge, therefore confinement leads to hadronisation of these outgoing partons, resulting in a shower of hadrons around the quark or gluon, as described in section 1.1.2. Instead of directly measuring the quarks, the inner detector and calorimeters detect the spray of particles, called a jet. The energy of the particles in the jet together provide an estimate of the energy of the outgoing quark or gluon. A jet reconstruction algorithm is needed to define which particles should be grouped together to form a jet.

Within ATLAS, only calorimeter information is used to reconstruct jets. The parton shower constituting a jet usually deposit their energy in many cells of the calorimeters. A topological clustering algorithm is used to sum the deposited energy of groups of cells, reconstructing the showers inside the calorimeter corresponding to each incoming particle. This algorithm starts with a seed cell with a high energy deposit, on top of which iteratively neighbouring cells are added if the energy in the new cell exceeds a given low energy threshold. If the energy of a neighbouring cell exceeds a higher threshold $t_{\text{neighbour}}$, it itself can be used as an additional seed, adding its neighbours to the list of potential cells.

These clusters are used as inputs for the jet finding algorithms, of which many have been defined. Ideally, these algorithms should be both infrared and collinear safe, meaning that emission of soft or collinear gluons by the outgoing partons should not affect the final reconstructed jets. In algorithms which are not infrared safe, soft gluon emission from an outgoing parton could increase the number of jets, which is unwanted. Collinear safety ensures that if collinear gluons redistribute their total transverse momentum over several particles, the p_T of the reconstructed jet is not affected.

The anti- k_t jet algorithm [185, 186], which is used in present ATLAS analyses, is indeed infrared and collinear safe, unlike the previously used fixed-cone algorithms. It produces conical jets using a sequential procedure based on the distance between the objects, and recombines the input objects (clusters) into jets until stable jets are found. The procedure is to calculate the distance measure d_{ij} between two objects i and j , and d_{iB} between object i and the beam, which are defined as

$$d_{ij} = \min(p_{T,i}^{-2}, p_{T,j}^{-2}) \frac{\Delta R_{ij}^2}{R} \quad (4.1)$$

$$d_{iB} = \frac{1}{p_{T,i}}, \quad (4.2)$$

where $p_{T,i}$ is the transverse momentum of object i and the distance parameter ΔR_{ij}^2 is defined as $\Delta R_{ij}^2 = \Delta\phi_{ij}^2 + \Delta\eta_{ij}^2$. The parameter R sets the size of the jet, and is chosen to be 0.4. For each object i a test is performed: if the difference d_{ij} is smaller than d_{iB} , objects i and j are recombined into a new object k , and the procedure is

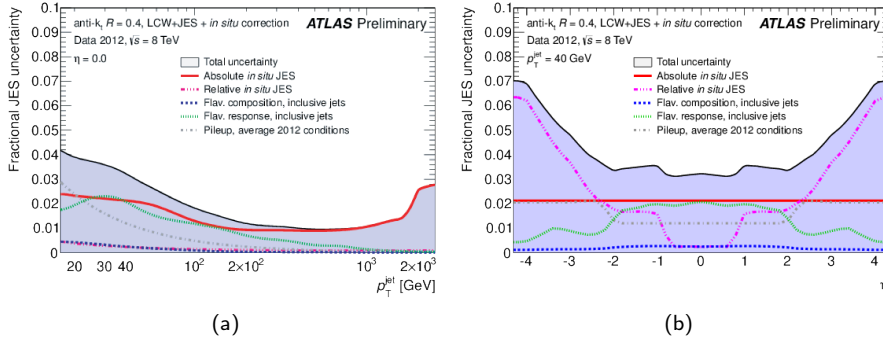


Figure 4.4: The jet energy scale uncertainty measured in 2012 as a function of (a) p_T and (b) η .

started anew. Once $d_{iB} < d_{ij}$ for all objects j , the object i is defined to be a jet and is removed from the list before searching for other jets. Due to the factor $1/p_{Ti}$, the clustering of objects into a jet starts with the object with highest p_T . Therefore, soft objects will not change the shape of the jet, and the jet algorithm is indeed infrared safe.

4.2.1 Energy scale

Since ATLAS uses sampling calorimeters, most of the energy of a jet is deposited in the absorbers. Other sources of reduced calorimetric response are leakage, jets which fall partially outside the calorimeter acceptance, and inefficiencies in jet clustering and reconstruction. Furthermore, the hadronic response is smaller than the electromagnetic response. To relate the measured energy by both calorimeters to the true energy of the shower of particles corresponding to the jet, they need to be calibrated to a jet energy scale (JES) [187]. The calibration scheme used in this thesis is the *Local Cluster Weighting* method (LCW+JES). As the name already suggests, this method calibrates the clusters locally before jet reconstruction. Each energy cluster is categorised as either electromagnetic or hadronic of origin, based on the shape of the cluster [188]. Due to this categorisation, an energy correction can be applied on each cluster corresponding to its (non)-electromagnetic nature. The energy corrections are derived from single pion MC simulation, with a dedicated correction for each of the above mentioned causes of differences between measured and true energy of a jet. From these locally calibrated clusters, jets can be formed using the previously described method. The LCW+JES procedure has lower JES corrections and a better jet energy resolution when comparing to other calibration schemes, such as EM+JES. For the latter, all clusters are first calibrated as being purely electromagnetic, upon which an additional factor is applied to the jet to correct for the lower hadronic response.

The systematic uncertainty on the jet energy scale is derived from *in situ* techniques from 2010 and 2011 collision data [187, 189], which has been checked to be consistent

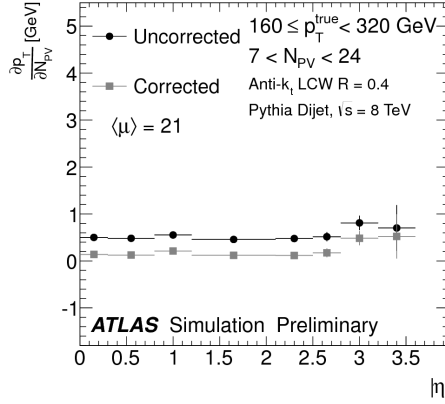


Figure 4.5: Dependence of the jet p_T on the number of primary vertices N_{PV} , representing the in-time pile-up dependence, as a function of $|\eta|$ measured in MC simulations at $\sqrt{s} = 8$ TeV. Here $\langle\mu\rangle = 21$, and the dependence is shown for uncorrected and pile-up corrected data. Taken from [192].

with the higher pile-up conditions in 2012. For central jets ($\eta = 0$ with $p_T < 2$ TeV, the uncertainty is less than 4%, while it is maximally 7% for very forward jets ($|\eta| \gtrsim 4$) with $p_T = 40$ GeV, as seen in figure 4.4. The systematic uncertainty on the jet energy resolution is obtained by comparing measurements in data and MC simulation of the di-jet balance and a bisector technique, which both are based on the momentum conservation in the transverse plane for a two jet event [190,191].

Jet reconstruction is affected by pile-up, due to the fact that jets are spread out over a wide area in the calorimeters, which increases the chance of particles originating from pile-up interactions to fall in the same area. The dependence of the jet response on these pile-up effects is decreased by introducing correction factors, derived from $\sqrt{s} = 7$ TeV MC simulation as a function of the average number of interactions per event and number of primary vertices. This is seen from the dependence of the average jet p_T on the number of primary vertices N_{PV} , shown in figure 4.5, where the red data points denote those with the pile-up correction applied. The data points are taken from the slopes of linear fits to distributions of $\langle p_T^{reco} - p_T^{true} \rangle$ versus N_{PV} , at a fixed $\langle\mu\rangle = 21$ value. Less dependence on pile-up, means a lower value of $\partial p_T / \partial N_{PV}$.

4.2.2 b -jets

Jets originating from b -quarks can be distinguished from jets coming from lighter quarks by using vertex information. As B -hadrons have a relatively long lifetime (~ 1.5 ps), they will travel typically several millimetres² from the interaction point before decaying and forming b -jets. These b -jets are crucial for the identification of the top quark backgrounds in the hadronic search for SUSY.

²Although $c\tau_B \sim 450$ μm , the flight distance of B -hadrons with a significant boost γ is enhanced up to several millimetres.

Many algorithms have been defined to tag b -jets, which all depend on the tracking and vertex information. These are described in ref. [193]. In this thesis, the MV1 algorithm has been used, which is based on a neural network which uses output of other algorithms as input – to be precise, of the IP3D, SV1 and JetFitterCombNN b -tagging algorithms. The output of the neural network is a distribution, upon which an operating point can be chosen which defines the exact requirements on the jet, from low to high tagging efficiency. A tagging algorithm should be efficient in the tagged variable, but also have a very low mistag rate: a very high efficiency in rejecting non- b -jets. The efficiencies have been derived from a dataset with jets and a muon using the p_T^{rel} method. p_T^{rel} is the muon momentum transverse to the axis of the combined muon plus jet system. Muons originating from a b -jet have a harder p_T^{rel} spectrum than muons in c - and light-flavour jets. This spectrum is fitted to multiple templates for b -, c - and light-flavour jets, to obtain the efficiencies in data, and scale factors to be used on MC simulation. Figure 4.6 (a) shows the tagging efficiency against the rejection of light-flavour jets (coming from MC generated up, down and strange quarks). In this thesis the 60% efficiency operating point was chosen, corresponding to a rejection factor of 635 and 8 for light-flavour jets and c -jets, respectively. In figure 4.6 (b) the b -tagging efficiency measured in a sample of $t\bar{t}$ events containing one muon is compared to MC simulation, using the 70% working point. Systematic uncertainties on the data to simulation scale factors are mostly due to uncertainties in heavy-flavour modelling of the MC generators, jet energy scale and resolution effects, and from limited MC statistics. The systematic uncertainties on the scale factors range from 5% and 20% for increasing jet p_T .

Note that due to the large decay width of the top quark (~ 1 GeV), corresponding to a lifetime of only 5×10^{-25} s [19], it decays before it can hadronise. Jets therefore cannot be associated to t quarks directly, only to their decay products.

4.2.3 Jet selection

Within the analysis described in chapter 5, jets are selected when they obey $p_T > 20$ GeV and $|\eta| < 2.8$. Jets are only identified as b -jets if they have $p_T > 40$ GeV and $|\eta| < 2.5$, on top of a b -tag described before. The rejection of badly reconstructed jets and non-collision backgrounds is described in section 5.3.1.

4.3 Leptons

Although the hadronic search for SUSY presented in chapter 5 targets jets, leptons are used, either to veto upon, or to select a sample of events from vector bosons or top quark pairs for control region samples. Therefore, the reconstruction and identification of leptons is of utmost importance. In the remainder of this thesis, with ‘leptons’ only electrons and muons are meant. Neutrinos are only identified by missing energy. Tau leptons decaying leptonically ($\tau \rightarrow l\bar{\nu}_l\nu_\tau$, with $l = e, \mu$) cannot be distinguished from electrons or muons, while hadronic tau decays end up as jets in this analysis.

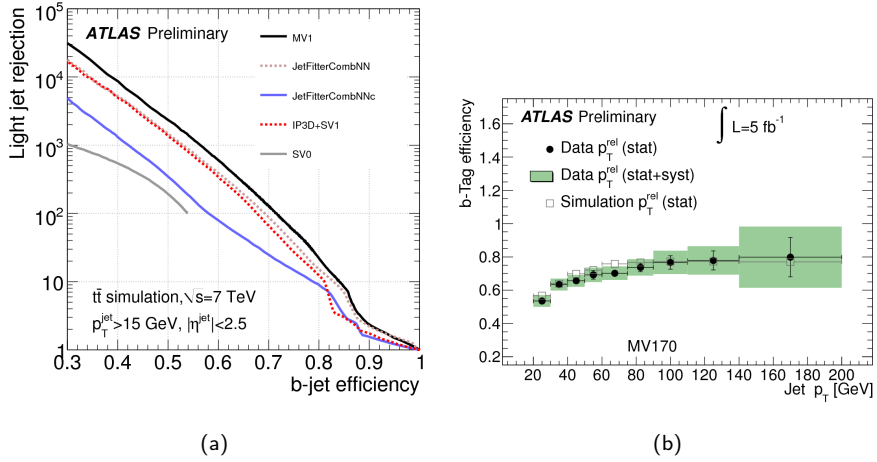


Figure 4.6: (a) The b -tagging efficiency as a function of the light-flavour jet rejection factor for various b -tagging algorithms, measured in $t\bar{t}$ simulation. (b) The measured b -tagging efficiency in a 5 fb⁻¹ $t\bar{t}$ sample containing one muon recorded in 2011 at $\sqrt{s} = 7$ TeV, compared to MC simulation. Taken from [193].

4.3.1 Muons

Muons are the only particles coming from the interaction point which survive the calorimeter (except for neutrinos), and are therefore mainly reconstructed and identified by the muon spectrometer, with additional information from the inner detector and calorimeters. Depending on which sub-detectors are used, various types of reconstructed muons can be identified: *stand-alone* muons are formed from information of just the muon spectrometer, where tracks are reconstructed and then extrapolated to the interaction point; *combined* muons are formed by reconstructing tracks in inner detector and muon spectrometer independently, after which a successful matching of an ID and an MS track is performed; finally *segment-tagged* muons are found by extrapolating inner detector tracks to the muon spectrometer, where they are matched to at least one straight track segment in the MDTs or CSCs. For each of these reconstruction types two separate algorithms have been defined, which are subsequently chained in two so-called muon collections, Chain 1 [194] and Chain 2 [195]. In this thesis, muons are reconstructed using the Chain 1 collection. While a fourth muon type exists, which uses muons tagged by the calorimeter and inner detector, it is only used in performance studies and thus not described here.

Stand-alone muons

The algorithms used to find stand-alone muons start by building straight track segments in each of the three muon stations near regions of activity identified by the

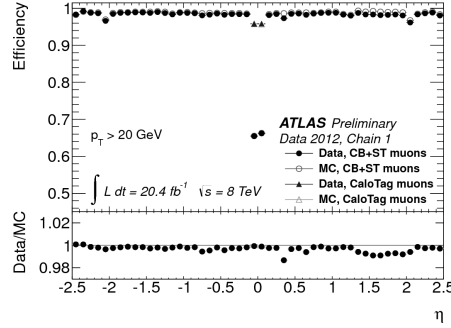


Figure 4.7: Reconstruction efficiency of muons as a function of η for different reconstruction types of the Chain 1 collection, in 20.4 fb^{-1} of data recorded in 2012. The lower pad shows the data to MC ratio. Taken from [196].

muon trigger chambers. The segments are built by combining hits in these stations and fitting a straight track through them. A stand-alone muon candidate track is reconstructed by a global fit of these segments extrapolated to the interaction region, while taking energy losses in the calorimeter and from multiple scattering into account.

The pseudorapidity coverage of stand-alone muons is determined by the muon spectrometer, which spans the range up to $|\eta| < 2.7$. However some chambers are missing at $\eta \sim 0$ to provide room for services, while at $\eta \sim 1.2$ some chambers have not yet been installed.

Combined muons

To combine the tracks found in the inner detector and muon spectrometer, both muon reconstruction chains perform a χ^2 fit of the matching of the inner and outer track vectors, using their combined covariance matrix. Here the inner detector tracks are found using the inside-out algorithm described in section 4.1. Chain 1 performs a statistical combination of inner and outer tracks, by extrapolating the inner track, combining it with the outer track closest to the extrapolation. Chain 2 on the other hand performs a refit of the track segments to achieve the best global fit of muon track, accounting for the magnetic field and material in front of the muon spectrometer.

Combined muons have the highest purity, yet due to the fiducial range of the inner detector, combined muons are restricted to $|\eta| < 2.5$.

Segment-tagged muons

To be able to reconstruct low- p_T muons which do not penetrate the whole muon spectrometer, segment-tagged muons are added to the collections. Here the inner detector tracks are extrapolated to the first muon station, where a search is performed for muon segments close to the predicted track position, based on either a χ^2 of the distance between the two (Chain 1) or a neural network (Chain 2).

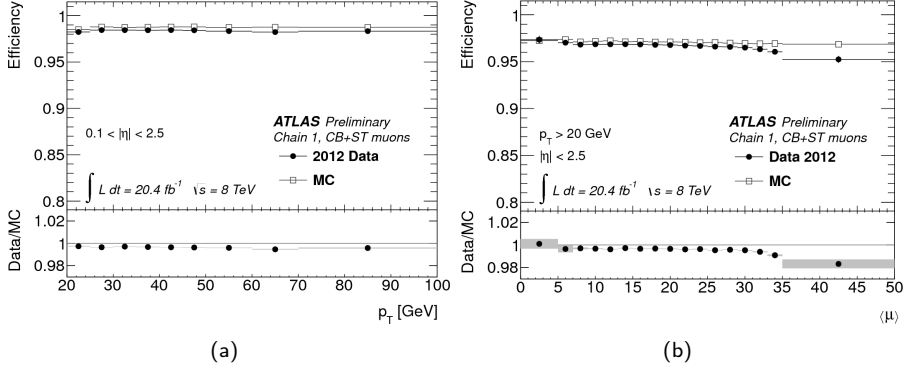


Figure 4.8: The reconstruction efficiency of muons using the Chain 1 collection as a function of (a) p_T and (b) the average number of interactions per event. The 20.4 fb^{-1} of data was recorded in 2012. The lower pad shows the data to MC ratio, with systematic uncertainties shown as a green band. Taken from [196].

Reconstruction efficiency and fake rate

The muon reconstruction efficiency is studied by a tag-and-probe method in a $Z \rightarrow \mu\mu$ sample: events are required to have two oppositely charged isolated muons with an invariant mass corresponding to the Z boson mass. Here one muon is tagged as a combined muon, while the other muon (the 'probe') is a calorimeter muon. Using these, the probability of reconstructing the second muon as a combined muons is studied. Figure 4.7 shows the efficiency for Chain 1 collection, for muons tagged as either combined or stand-alone muon (circles) or tagged by the calorimeter (triangles), taken from 2012 data and MC simulation generated with POWHEG [197]. An efficiency of $\sim 98\%$ is achieved everywhere except at $\eta \approx 0$ for combined and stand-alone muons, with slightly higher efficiency measured in MC.

Figure 4.8 (a) shows the efficiency of combined and stand-alone muons from Chain 1 as a function of the transverse momentum of the muons. The efficiency is seen to be nearly independent of p_T . The pile-up dependency of the efficiency is shown in figure 4.8 (b), which is quite stable as well, with only a small drop in efficiency for a very large amount of pile-up, at $\langle \mu \rangle$ greater than 35.

The difference in reconstruction efficiency seen in MC and data leads to the introduction of a scale factor to be applied to MC, correcting the efficiency behaviour to reflect the behaviour in data. It is defined as just the η and p_T dependent ratio of the efficiencies ϵ_{data} and ϵ_{MC} in data and MC, respectively:

$$SF = \frac{\epsilon_{data}}{\epsilon_{MC}}. \quad (4.3)$$

For the muon veto used in the hadronic SUSY analysis described in chapter 5,

the number of ‘fake’ muons should be small, as it would lead to events being falsely rejected. These so-called ‘fake’ muons originate mainly from semileptonic heavy-flavour decays inside jets; they are therefore real muons, yet do not come from the prompt decay. The fake rate, or relative number of events which contain such a fake muon, is calculated using MC simulation: for a SUSY signal model without leptons in the decay, the number of events before and after a muon veto is compared. For a squark-pair model with $m_{\tilde{q}} = 450$ GeV, where the produced squarks each decay directly to a quark and an LSP, and with 60000 simulated events, the obtained muon fake rate is $0.09\% \pm 0.01\%$, where the statistical uncertainty is given by the binomial error $\Delta\epsilon = \sqrt{\epsilon(1-\epsilon)/N}$, with ϵ the fake rate and N the number of MC events.

Further muon selection

To select muons for physics analyses, the muons need to pass quality requirements on the number of hits in each inner detector component: at least 1 pixel hit and B -layer hit (when expected); at least 6 SCT hits; less than 3 holes in total in the SCT and pixel detectors; and at least 6 TRT hits for $|\eta| \leq 1.9$, while for $|\eta| > 1.9$ no TRT requirement is applied. Of these TRT hits less than 90% may be outliers: hits associated to the track but which result in a bad combined tracking fit, most probably due to an unsuccessful extrapolation from ID tracks to the TRT. Selected muon tracks should be either reconstructed as a combined or as a segment-tagged muon. A so-called ‘baseline’ muon, used to veto muons, should have a transverse momentum of more than 10 GeV, and it should fall in $|\eta| < 2.4$. ‘Signal’ muons, used for muon selections in control regions, are further required to be isolated, by ensuring that the sum of the p_T of all charged particle tracks, associated with the primary vertex and which are within $\Delta R \equiv \sqrt{(\Delta\phi)^2 + (\Delta\eta)^2} < 0.2$ of the muon, is less than 1.8 GeV. To ensure high reconstruction efficiency with high purity, leading signal muons are required to have $p_T > 25$ GeV, while a second signal muon should have $p_T > 20$ GeV. To reject muons coming from outside the detector (for instance from cosmic rays), the impact parameters of the reconstructed muon should obey $|z_\mu - z_{PV}| < 1$ mm and $d_0 < 0.2$ mm. Finally, the muon momentum is smeared in MC simulation to obtain an equal dimuon mass resolution to that measured in data.

Figure 4.9 (a) shows the dimuon invariant mass spectrum, obtained from 40 pb^{-1} of $\sqrt{s} = 7$ TeV data, where the events are selected by the event filter on information from both the ID and MS. In the shown reconstructed dimuon invariant mass distribution several known resonances can be found, with the Z boson on the right at a mass of 91 GeV. Figure 4.9 (b) shows the mass resolution of this peak in several pseudorapidity ranges using 205 pb^{-1} of 2011 data, and comparing to Monte Carlo simulation. The mass resolution is observed to be between 3 and 5 GeV.

4.3.2 Electrons

Electrons leave a track in the inner detector, and are stopped in the electromagnetic calorimeter. Therefore information of both sub-detectors is used to reconstruct electrons. Reconstruction starts by clustering calorimeter cells using a *sliding window*

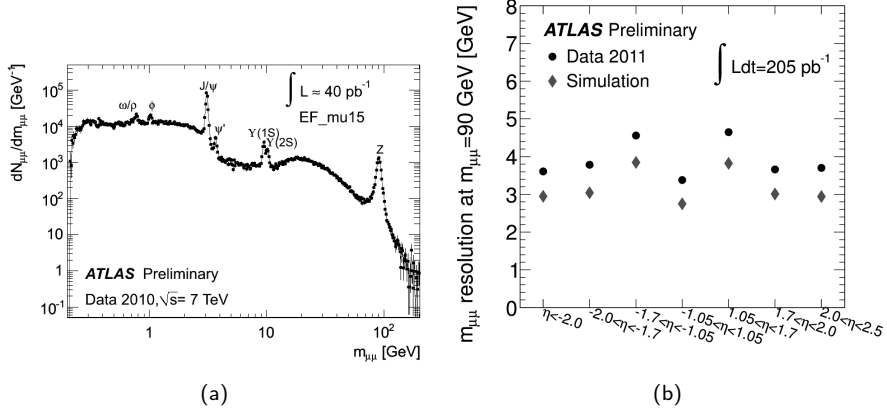


Figure 4.9: (a) The dimuon invariant mass spectrum using 40 pb^{-1} $\sqrt{s} = 7 \text{ TeV}$ data. (b) The mass resolution of the $m_{\mu\mu}$ invariant mass, made with 205 pb^{-1} $\sqrt{s} = 7 \text{ TeV}$ data from the muon spectrometer alone, in bins of η .

algorithm [198]. Here a rectangular window the size of 15 middle layer cells, or 3×5 units of 0.025×0.025 in $\eta \times \phi$ space, is moved over the calorimeter cells, until its summed transverse energy exceeds a threshold value of 2.5 GeV. The cluster corresponding to the local maximum is used as a seed cluster, which is matched to an inner detector track. Tracks are extrapolated from their last measurement point to the EM calorimeter, where they should be within $\Delta\eta < 0.05$ of the seed cluster. Due to Bremsstrahlung, electrons lose energy and thus their trajectories may change substantially in the ϕ direction. These losses are taken into account using the Gaussian Sum Filter (GSF) [199], which models the losses as a sum of Gaussian functions. An electron is reconstructed if at least one track is matched to the seed cluster. In case of multiple tracks pointing to the cluster, preference is given to tracks with SCT hits, while finally the track with smallest ΔR is chosen. The cluster energy is defined by summing the contributions in the EM calorimeter together with the estimated energy which is lost, either due to material in front of the calorimeter, energy missed during clustering, or energy deposited beyond the calorimeter. The electron four-momentum is computed from a combination of the calorimeter information and the tracks refitted by the GSF: the energy is taken as the cluster energy, while the η and ϕ parameters are taken from the track information if the track has at least 4 silicon hits. If it does not, η is taken from the EM cluster. Clusters are only taken into account if they can have a track matched to them, thus they should be inside the inner detector acceptance, $|\eta| < 2.47$.

Many electron candidates reconstructed with the procedure above will have originated from photon conversions or from jets faking an electron. Real electrons are identified by using requirements on both the tracks and the calorimeter information.

Three η - and E_T -dependent selections are defined, with increasing background rejection: *loose++*, *medium++* and *tight++*. For a loose selection, only the shower shape information from the middle EM calorimeter layer is used. The medium selection requires a matched track which passes quality requirements: at least 7 hits in the pixel and SCT combined, of which at least one should be in a pixel layer; the matched track should be close in η to the cluster, $|\eta_{cluster}^{1st\,layer} - \eta_{track}| \leq 0.01$; and the distance d_0 to the primary vertex should be small, $d_0 \leq 5$ mm. It also utilises the shower shape in the first layer of the calorimeter. This selection is tightened even more in the tight selection by adding a requirement on the ratio of the electron energy over its momentum, E/p , and particle identification using transition radiation in the TRT. The track should be even closer to the cluster, $\Delta\phi \leq 0.02$ and $\Delta\eta \leq 0.005$. To protect against photon conversions, a B-layer hit and a veto on a reconstructed photon vertex is required as well. The postfix ‘++’ in the purity selections comes from modifications introduced in 2012 to reduce the effect of pile-up on the selection.

‘Baseline’ electrons are defined as electron candidates passing medium++ requirements, with $|\eta| < 2.47$ and $p_T > 20$ GeV. ‘Signal’ electrons, used in control regions requiring at least one lepton, are required to pass the tight++ cuts, and have $p_T > 25$ GeV. Furthermore, to protect against photon conversions, the transverse and longitudinal distance to the primary vertex should be less than 1 mm and 2 mm, respectively. Finally, to select an isolated electron, the sum of the p_T of all charged particle tracks associated with the primary vertex within $\Delta R < 0.2$ of the electron should be less than 10% of the electron p_T .

The reconstruction efficiency is measured using a tag-and-probe method in a sample of $Z \rightarrow e^-e^+$ events, selected by requiring two oppositely charged isolated electrons with an invariant mass corresponding to the Z boson mass, of which one is a ‘tight’ electron while the other selected using looser cuts. The efficiency is calculated from the number of times the ‘probe’ is reconstructed using the ‘tight’ criteria. Figures 4.10 (a) and (b) show the reconstruction efficiency of electrons as a function of η and E_T . The improvement in 2012 due to the introduction of the Gaussian Sum Filter is clearly visible. The identification efficiency, using the loose, medium and tight selections, is shown in figure 4.10 (c) as a function of the number of primary vertices. The flatness of these distributions shows the robustness against pile-up effects.

The differences between data and the MC simulation in the reconstruction efficiencies, and also in the reconstructed isolation energy, leads to η dependent scale factors which are applied to electrons identified in data. Uncertainties on the electron reconstruction come from the reconstruction efficiency, and the electron energy scale and resolution, where the efficiency uncertainty is measured through tag-and-probe methods.

‘Fake’ electrons originate mainly from charged hadrons misidentified as electrons, photon conversions or semileptonic heavy-flavour decays inside jets. The electron fake rate, calculated similarly to the muon fake rate in hadronically decaying SUSY signal samples, is $0.20\% \pm 0.02\%$.

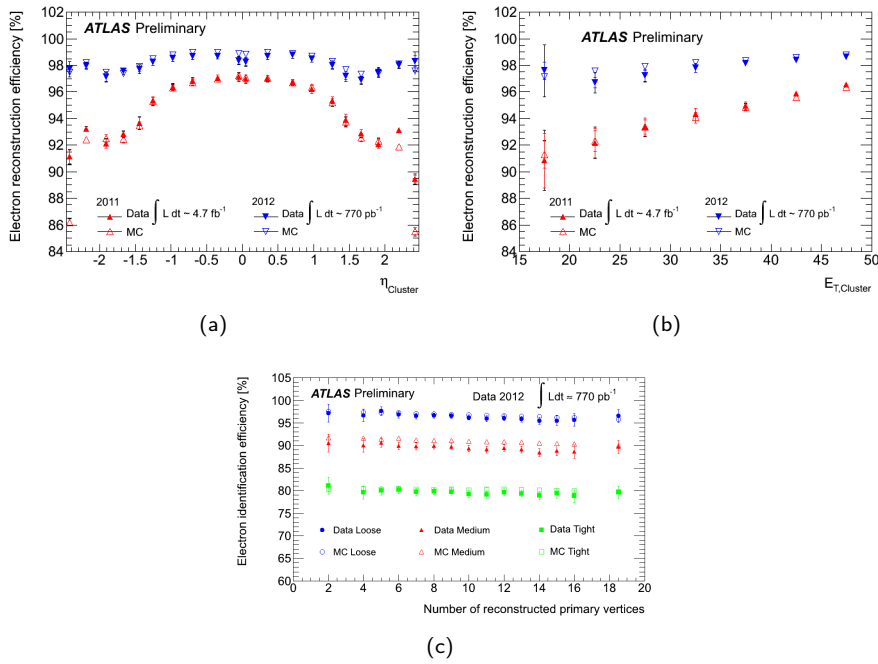


Figure 4.10: The reconstruction efficiency of electrons as a function of (a) η and (b) E_T , in 2011 (4.7 fb^{-1}) and 2012 (770 pb^{-1}), and (c) as a function of the number of primary vertices in 770 pb^{-1} of $\sqrt{s} = 8 \text{ TeV}$ data, for the three different purity selections [200]. The error bars show the total statistical and systematic uncertainty on each measurement.

4.4 Photons

Photons are only used in one control region of the analysis presented in chapter 5. Their reconstruction in the calorimeter is similar to that of electrons. Photons can either enter the EM calorimeter unconverted, or can convert into an e^-e^+ pair in the inner detector. While the former do not leave any track, the latter are recovered by matching the tracks to the electromagnetic clusters. The EM clusters are identified using the same sliding window algorithm as for electrons. Apart from the lack of a matched track, unconverted photons are differentiated from electrons by their shower shapes, which tend to be narrower for photons. A background for unconverted photons are neutral pion decays ($\pi^0 \rightarrow \gamma\gamma$), which can lead to two photons detected in the calorimeter. Using the high granularity of the first layer of the EM calorimeter, the two photons can usually be separated, leading to a π^0 reconstruction.

Photon candidates are required to have a transverse energy passing the trigger requirements of the lowest unprescaled photon trigger, $E_T > 130 \text{ GeV}$. However, photons

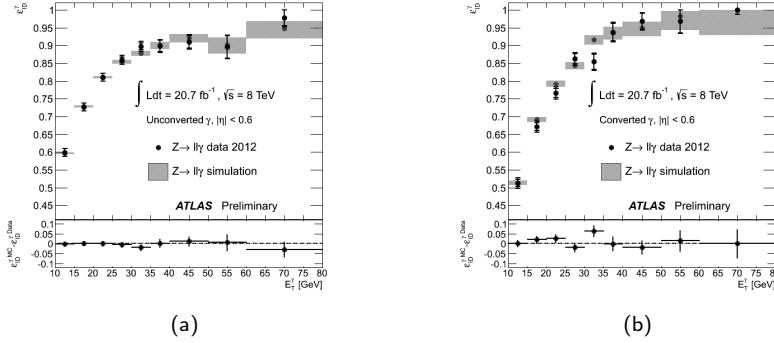


Figure 4.11: The identification efficiency of (a) unconverted and (b) converted photons for $|\eta| < 0.6$, as a function of the photon E_T [201].

used for $\cancel{E}_T$ reconstruction (see section 4.5) are only required to have $E_T > 10$ GeV. Photon candidates are required to be in the fiducial region of the EM calorimeter, $|\eta| < 1.37$ and $1.52 < |\eta| < 2.37$, as photons in the transition area have a significantly lower reconstruction efficiency. Isolated photons are selected by requiring that the energy in a cone of $\Delta R < 0.4$ around the photon candidate is less than 4 GeV, when removing the photon energy.

The identification efficiency of photons is measured from radiative Z decays, $Z \rightarrow l^+l^-\gamma$ with $l = e, \mu$. This sample is obtained by requiring the reconstructed dilepton and dilepton plus photon mass to obey $40 \text{ GeV} < m_{ll} < 83 \text{ GeV}$ and $40 \text{ GeV} < m_{ll\gamma} < 96 \text{ GeV}$, respectively, with $E_T^\gamma > 10 \text{ GeV}$, while no cuts are placed on the photon shower shape variables to prevent a bias on the efficiency measurement. The efficiency is shown in figure 4.11 for 20.7 fb^{-1} of $\sqrt{s} = 8 \text{ TeV}$ data, for the central region. The efficiency increases for higher E_T , with an efficiency of around 90% for photons with $p_T > 35 \text{ GeV}$. As the photon spectrum falls steeply with its E_T , the uncertainties for high E_T are statistically driven and become very large for $E_T > 50 \text{ GeV}$.

4.5 Missing transverse energy

So far, all the described reconstructed physics objects were of detectable particles. Yet neutrinos, and also LSPs in the case of an R-parity conserving SUSY signal, do not interact with the detector, and therefore leave no detectable trace. Some information on the presence of these particles can be obtained however, using the momentum imbalance in the transverse plane. As the initial state has no momentum in the transverse plane, the vector sum of the momentum of all final state particles should not have any transverse component due to momentum conservation. However, the measured net transverse momentum will be non-zero if any final state particle was not measured, as would be the case for a neutrino or LSP. The measured missing

transverse momentum is thus equal to the net transverse momentum of the invisible particles, $\vec{p}_T^{\text{miss}} = \sum_{\text{inv. particles}} \vec{p}_T$. On the other hand, it is calculated as the negative net transverse momentum of all visible particles, $\vec{p}_T^{\text{miss}} = -\sum_{\text{vis. particles}} \vec{p}_T$. Assuming the invisible particles are massless, we take the magnitude of the missing transverse momentum as $p_T^{\text{miss}} = \cancel{E}_T = \sqrt{\cancel{E}_x^2 + \cancel{E}_y^2}$.

However, many effects can pollute the $\cancel{E}_T$ measurement. Misreconstructed objects, calorimeter noise, and pile-up events can all give rise to energy mismeasurements. Furthermore, as ATLAS has no full 4π coverage, particles can escape undetected. Therefore many systematic uncertainties need to be taken into account when reconstructing $\cancel{E}_T$.

The reconstruction of $\cancel{E}_T$ is performed by combining information of all physics objects described in the previous sections, using mainly the energy deposits in the calorimeters, together with muon information. Calorimeter clusters associated to electrons, photons, and jets are used as part of the physics objects, instead of using the clusters themselves, so as to be able to calibrate each cluster separately based on the object it belongs to. This leads to components of the $\cancel{E}_T$ for each of these objects: $\cancel{E}_T^{\text{ele}}$, $\cancel{E}_T^{\text{photon}}$ and $\cancel{E}_T^{\text{jet}}$, which are all taken as the *negative* sum of the cluster energies due to momentum conservation. The electron contribution comes from ‘baseline’ electrons, passing medium++ purity cuts with $p_T > 10$ GeV. The photon component is added for selected photons with $p_T > 10$ GeV. The contribution from jets is added for jets with $p_T > 20$ GeV after the LCW+JES calibration, while softer jets with $7 \leq p_T \leq 20$ GeV using the LCW calibration without jet energy scale corrections are added in a separate $\cancel{E}_T^{\text{SoftJets}}$ term. Clusters which do not belong to any of these objects are assumed to correspond to hadronic activity which has not formed jets, and they are gathered in a CellOut term $\cancel{E}_T^{\text{CellOut}}$ using the local LCW calibration.

Finally, although muons are not stopped in the calorimeter, they contribute as well to $\cancel{E}_T$. The muon term $\cancel{E}_T^{\text{muon}}$ is taken as the negative sum of the reconstructed p_T of the muon tracks for ‘baseline’ muons with $p_T > 10$ GeV, without any isolation requirement. Both electron and muon contributions are taken before the removal due to overlapping objects has been performed, as described in section 4.6. The missing transverse energy in the x and y direction is thus defined as

$$\cancel{E}_{x(y)} = \cancel{E}_{x(y)}^{\text{ele}} + \cancel{E}_{x(y)}^{\text{photon}} + \cancel{E}_{x(y)}^{\text{jet}} + \cancel{E}_{x(y)}^{\text{muon}} + \cancel{E}_{x(y)}^{\text{SoftJets}} + \cancel{E}_{x(y)}^{\text{CellOut}}. \quad (4.4)$$

To remove any ambiguity from clusters belonging to multiple objects, each of the components is added in the given order, and overlapping clusters are removed from later components. In this definition of $\cancel{E}_T$ the contribution of τ -decays is not explicitly included, yet hadronically decaying τ -leptons will be taken into account via the jet terms. The $\cancel{E}_T$ performance has been analysed in $Z \rightarrow l^+l^-$ and $W^\pm \rightarrow l^\pm\nu$ events on 20 fb^{-1} of $\sqrt{s} = 8$ TeV data in 2012 [202]. In the former sample, ‘fake’ $\cancel{E}_T$ coming from detector mismeasurements can be studied, while in the latter sample the neutrino leads to real $\cancel{E}_T$, which can be used to validate the $\cancel{E}_T$ scale. A comparison

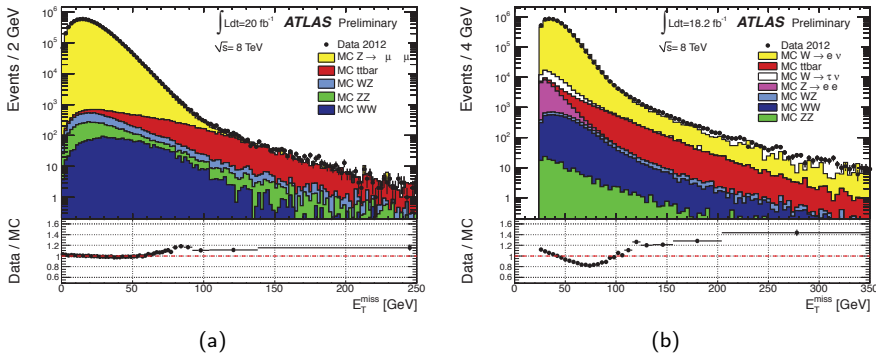


Figure 4.12: Distribution of E_T in 2012 $\sqrt{s} = 8$ TeV data in (a) $Z \rightarrow \mu\mu$ and (b) $W \rightarrow e\nu$ events, compared to MC simulation. The MC histograms are superimposed on top of each other. The lower parts show the ratio of data over MC. Figure taken from [202].

of the reconstructed E_T for these events in data and MC simulation generated with PYTHIA6 is shown in figure 4.12. Here the MC simulation for the Z or W production and their backgrounds are each normalised to their corresponding cross section and are superimposed. A good agreement is between the $Z \rightarrow \mu\mu$ data sample and MC simulation, while for $W \rightarrow e\nu$ events the low E_T region is not well described very well, with a discrepancy up to 40% for high E_T events. This is most likely due to the fake electron background, which is not included in the shown MC expectation [202].

The E_T resolution can be taken from $Z \rightarrow ll$ events, as these produce no real E_T . Using the width of the combined E_x and E_y distribution as a function of $\sum E_T$, the resolution can be approximated by a stochastic function of the total transverse energy, $\sigma = k \cdot \sqrt{\sum E_T}$ [203]. The thus obtained resolution is shown for $Z \rightarrow ee$ and $Z \rightarrow \mu\mu$ events in figure 4.13, with a fit using the approximate function given before. The fitted value for k of ~ 0.7 GeV^{1/2} has been shown to agree well with MC. Yet comparing to the observed resolution in 2010, with $k \sim 0.5$ GeV^{1/2}, it has degraded due to the increased pile-up conditions. However, the tight selection on jets and E_T used in the inclusive hadronic search for SUSY on 5.8 fb⁻¹, presented in this thesis, reduces the effect of pile-up significantly, as will be shown in chapter 5.

Uncertainties on the jet energy scale and resolution, as well as on the lepton scale and resolution are propagated into the systematic uncertainty on E_T . Furthermore, the contributions of the CellOut and SoftJet terms bring their own uncertainty, which are evaluated separately. These have been determined from *in situ* techniques on $Z \rightarrow \mu\mu$ events [204]. The total systematic uncertainty on E_T is evaluated on $W \rightarrow e\nu$ events and is shown to range from 2% to 10%, increasing with $\sum E_T$.

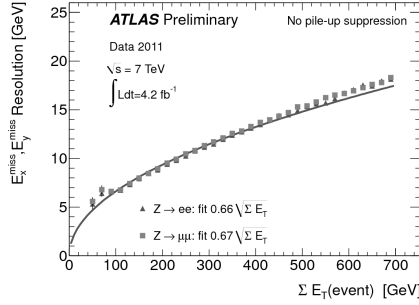


Figure 4.13: Distribution of the E_T resolution as a function of $\sum E_T$ for $Z \rightarrow ee$ (triangles) and $Z \rightarrow \mu\mu$ (squares) events. The fit is only shown for the $Z \rightarrow ee$ resolution. Figure taken from [204].

4.6 Overlapping objects

Physics objects overlapping in (η, ϕ) need to be treated with care. Clusters in the electromagnetic calorimeter can be used for the reconstruction of jets, while at the same time being identified as coming from electrons or photons. On the other hand, leptons close by the jet axis ($\Delta R < 0.4$) are more probable to originate from the decay of a heavy-flavour quark or τ in the jet, than from the initial interaction. Furthermore, jets can punch through the hadronic calorimeter, leading to a signal in the muon spectrometer. To resolve these issues, the removal of overlapping objects is performed using the geometric variable $\Delta R = \sqrt{\Delta\phi^2 + \Delta\eta^2}$ for any jet candidate with $p_T > 20$ GeV and any electron, muon or photon candidates in the following order:

- 1 If the distance between an electron and a jet is $\Delta R < 0.2$, the object is assumed to be an electron, and hence the jet is discarded;
- 2 If the distance between a muon and a jet is $\Delta R < 0.4$, the object is assumed to be a jet, and hence the muon is discarded;
- 3 If the distance between an electron and a jet is $0.2 < \Delta R < 0.4$, the object is assumed to be a jet, and hence the electron is discarded;
- 4 In case of photon selection: if the distance between a photon and a jet is $\Delta R < 0.2$, the object is assumed to be a photon, and hence the jet is discarded.

4.7 Monte Carlo simulation

In order to gain confidence in the understanding of a result and validate performance estimates, observations are compared to theoretical predictions – as has been done in this chapter for many performance results. In case of high energy collider physics, this means that a simulation of the collisions and subsequent processes needs to be

performed to be able to compare observed events with theoretical expectations. These simulations should also include a realistic simulation of the ATLAS detector to mimic its response. Simulations of Standard Model and SUSY signal processes are also used to design and optimise physics analyses and interpret results.

4.7.1 Event generation

Event generation is done using Monte Carlo (MC) generators: computational algorithms which rely on random numbers to simulate the stochastic nature of the underlying quantum field theories. The simulation of events consists of several steps, which are separated due to their different energy scales: first the *hard scattering* is calculated, after which the *parton showers* and *hadronisation* of the partons are performed together with the subsequent decays of outgoing particles.

Hard scattering

For each candidate event, the incoming partons are assigned a position and four-momentum by a random number generator, according to the PDFs. Using the factorisation theorem, the matrix elements belonging to the hard process are calculated perturbatively, while the PDFs describing the momentum distribution of the incoming protons are taken as input to the generators. Each generated event is scaled to correspond with the differential cross section, obtained from the matrix element, PDFs and kinematics. This leads to hard scattering simulation with the characteristics predicted by the Standard Model (or new physics).

Parton showers and hadronisation

Partons involved in the collision can radiate off gluons, resulting in either initial state radiation (ISR) in case of radiation off partons before the hard interaction, or final state radiation (FSR) off partons coming from the hard interaction. These gluons can subsequently split into quark-antiquark pairs or additional gluons, leading to a parton shower. Low energy and collinear radiation cannot be calculated perturbatively, and the radiation is thus implemented using parton splitting functions describing the possibility for a parton to split into two. Several MC generators can perform this parton showering, using the DGLAP [35–37] evolution equations and Sudakov form factors [205] to include the soft and collinear gluons.

The product of parton showers are still coloured partons. These should be transformed into colourless mesons and baryons, which can be observed. This hadronisation of outgoing partons into jets likewise cannot be calculated perturbatively, but is implemented in the generators using phenomenological models which are tuned to data to reproduce the observed jet properties in ATLAS and elsewhere as well as possible. The *Lund string fragmentation* [206] and *cluster fragmentation* [207] models are used by the PYTHIA and HERWIG MC generators, respectively. Possible hadron decay is treated after hadronisation using experimentally obtained branching ratios wherever available.

Matching parton showers to the matrix element

A problem arises when looking closer at the simulation of the hard scattering and the parton shower: a hard process with three outgoing partons cannot be kinematically discriminated from one with two outgoing partons, of which one radiates off a gluon via the parton shower. This may lead to double counting of events, one for each possibility, in particular in events with many jets [208]. Therefore a prescription (matching scheme) is needed to define which path, either using the hard process or parton showering, is used to generate the specific event.

Each of the schemes works along the same principles: the phase space available for parton emissions is divided into hard and large-angle emissions, handled by the matrix elements, and soft and collinear emissions, described by parton showers. The generators used in this thesis use either the MLM scheme [209] or the CKKW scheme [210–212]. The former, used in ALPGEN and PYTHIA, implements a veto on events which contain an emission from the parton shower in the matrix element region. The CKKW scheme, implemented in SHERPA, work similarly but with reweighted Sudakov form factors. Two next-to-leading order (NLO) generators, MC@NLO and POWHEG, use event weights as matrix element corrections to account for double counting of events. For more information, see ref. [44].

4.7.2 Generators

Due to the various implementations of the hard process calculation, the modelling of parton showering and hadronisation, and the differences in matching schemes, there are many MC generators used within the high energy physics community, of which several have already been introduced. The generators used for simulated data in this thesis are:

PYTHIA [142]: a leading order (LO) general purpose generator, able to simulate the hard scattering for processes with 2 incoming and 1 or 2 outgoing partons ($2 \rightarrow 2$), as well as the parton showering and hadronisation of these events using string fragmentation. Multi-parton final states are only obtained through the parton shower simulation, thus reducing the accuracy. PYTHIA is used for the QCD multi-jet simulations, and is the main method of parton showering for pure matrix element generators.

HERWIG [213, 214]: a LO general purpose generator similar to PYTHIA, yet using cluster fragmentation. HERWIG is used in combination with other generators, providing the parton showering. An improved version, HERWIG++ [215] is used for the simulation of several SUSY signals. HERWIG is interfaced with JIMMY [216] for the simulation of the underlying event, which is a MC generator dedicated to this specific task.

ALPGEN [217]: a leading order matrix element generator, which has as output outgoing partons. ALPGEN can generate up to six additional partons in the matrix element. Therefore it is mostly used for simulation of vector boson plus jets

events. To perform the parton showering for the outgoing partons, ALPGEN can be interfaced with PYTHIA or HERWIG.

SHERPA [43]: a LO matrix element generator similar to ALPGEN. However, SHERPA can perform the parton showering itself. It is again mostly used for vector boson generation with additional jets.

MC@NLO [218–221]: calculates the matrix element up to next-to-leading order. It is interfaced with HERWIG for the parton showering and hadronisation. It is mainly used for the simulation of top quark events.

POWHEG [197, 222–224]: another NLO matrix element generator, interfaced with PYTHIA for the parton showering and hadronisation. It is again mainly used for the simulation of top quark events.

ACERMC [225]: a LO generator dedicated to Standard Model background processes at the LHC, used in the analyses presented in this thesis for top quark production. Two different tunes of PYTHIA are used to evaluate the amount of additional ISR/FSR radiation, with a harder or softer parton shower.

MADGRAPH [226]: a general purpose matrix element based generator at LO, mostly used for SUSY signal production.

4.7.3 Simulation of the ATLAS detector

The MC generators provide a set of generated events which still consist of particles, comparable to a real collision inside the detector. To transform these into detector signals, they need to be propagated through a simulation of the ATLAS detector. This is done in three stages: the detector simulation, digitisation and event reconstruction. In the first, interactions of the simulated particles with the detector are mimicked using GEANT4 [227]. During digitisation, the detector response to these interactions is modelled, giving an identical output to real particles traversing the detector [228]. Finally, event reconstruction is performed identically to real data, as described previously.

Search for Supersymmetry in events with jets and $\cancel{E}_T$ in 5.8 fb^{-1} of $\sqrt{s} = 8 \text{ TeV}$ data

In the search for supersymmetry, many different signatures are probed. In this thesis we concentrate on SUSY signatures containing at least two jets with large momentum together with large missing transverse momentum, without any light leptons (electrons or muons): the so-called ‘0-lepton’ analysis.¹ The described analysis is published in the ATLAS note [172] using an integrated luminosity of 5.8 fb^{-1} at $\sqrt{s} = 8 \text{ TeV}$, which itself is a continuation of the analysis on 4.7 fb^{-1} at $\sqrt{s} = 7 \text{ TeV}$ presented in ref. [169].

Like most large analyses conducted at ATLAS, this analysis was performed by a collaboration of physicists. Amongst the ~ 10 collaborators working closely on the analysis, I have concentrated most on the analysis for compressed SUSY scenarios. For the compressed scenarios I have updated the signal grids for direct squark-pair and gluino-pair production and introduced a new squark-gluino pair-production grid, performed analysis optimisation for these grids from start to finish, leading to new Signal Regions and setting exclusion limits from the final results. To improve the background estimation I have worked on the comparisons between Monte Carlo generators for the various backgrounds and defining final selections within the Control Regions for $t\bar{t}$, W +jets and Z +jets backgrounds. Furthermore, I have worked on analysis framework development and maintenance.

This chapter first introduces shortly the targeted SUSY signals, their backgrounds and the followed analysis strategy. The dataset and Monte Carlo samples used are described in section 5.2, after which the event selection and optimisation procedure used to define the signal regions is presented in section 5.3. Section 5.4 describes the background estimation procedure, while the statistical procedure for used in both the background estimation and quantifying the results is described in section 5.5. Finally

¹Although neutrinos are leptons as well, in this chapter ‘lepton’ is only used to describe electrons and muons, and tau-leptons where specifically stated.

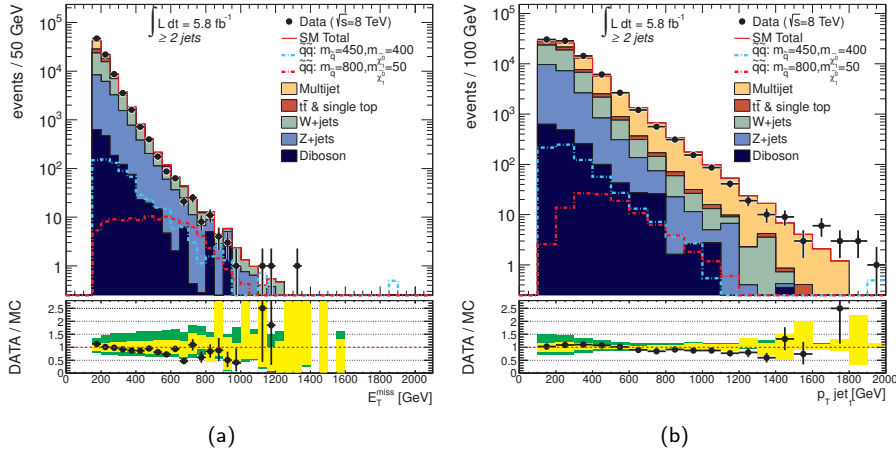


Figure 5.1: Observed and simulated (a) $\cancel{E}_T$ and (b) p_T of the leading jet in events with no electrons or muons, $\cancel{E}_T > 160$ GeV and at least 2 jets (with $p_T > 130$ and 60 GeV, respectively) in 5.8 fb^{-1} of $\sqrt{s} = 8$ TeV data. The dashed histograms show the distribution of two signal models, where squark-pairs are produced and decay directly, with $m_{\tilde{q}} = 450$ GeV, $m_{\tilde{\chi}_1^0} = 400$ GeV (light blue) and $m_{\tilde{q}} = 800$ GeV, $m_{\tilde{\chi}_1^0} = 50$ GeV (red). In the bottom pad the data to MC ratio is given, with the combined systematic uncertainties in illustrated by the yellow band, and the total uncertainty including theory uncertainties in the green band.

the results are given in the last section. The discussion of these results is left for the next chapter.

5.1 Overview of the 0-lepton analysis strategy

If SUSY exists, the most abundantly produced SUSY particles at the LHC will be squarks and gluinos, provided that they are not too heavy, due to the coupling strength of the strong force, as explained in section 1.3.4. These squarks and gluinos decay either directly into quarks forming jets, or via cascade decays resulting in jets and possibly leptons or photons. In R-parity conserving SUSY, the LSP travels through the detector undetected, leading to missing transverse momentum p_T^{miss} , the magnitude of which we denote as $\cancel{E}_T$.

Of the many available SUSY models, both the CMSSM and three simplified models are targeted. The used simplified models focus on the $\tilde{q}\tilde{q}$, $\tilde{g}\tilde{g}$ and $\tilde{q}\tilde{g}$ production and decay directly into quarks and LSPs, and are presented in section 1.3.4. Since within the CMSSM $\tilde{g}\tilde{g}$, $\tilde{q}\tilde{q}$, $\tilde{q}\tilde{q}$ and $\tilde{q}\tilde{g}$ are all produced, for all models under consideration we have the following signatures:

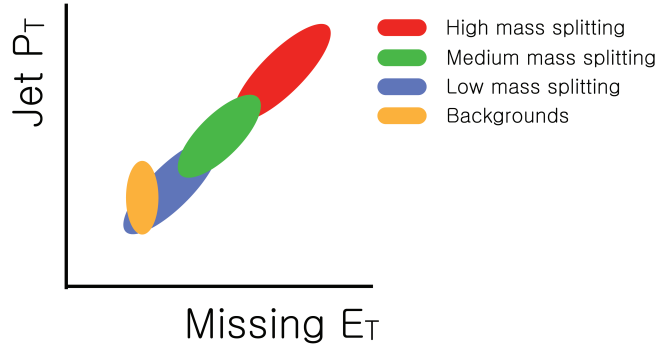


Figure 5.2: Schematic plot of the typical $\cancel{E}_T$ and leading jet p_T for directly decaying SUSY scenarios with low (light blue), medium (green) and high mass splitting (red), defined as the mass difference between the strongly interacting SUSY particle and the LSP. The typical values for QCD, Top, Z+jets and W+jets backgrounds is shown in orange.

- $\tilde{q}\tilde{q}$ or $\tilde{q}\tilde{\bar{q}}$: Each squark decays into a quark and a $\tilde{\chi}_1^0$, resulting in 2 jets plus $\cancel{E}_T$.
- $\tilde{q}\tilde{g}$: The squark decays into a quark and a $\tilde{\chi}_1^0$, while the gluino decays into 2 quarks and a $\tilde{\chi}_1^0$, resulting in 3 jets plus $\cancel{E}_T$.
- $\tilde{g}\tilde{g}$: Each gluino decays into 2 quarks and a $\tilde{\chi}_1^0$, resulting in 4 jets plus $\cancel{E}_T$.

Cascade decays, i.e. longer decay chains, are possible as well within the CMSSM, while additional jets from initial and final state radiation (ISR and FSR) are obviously possible in all events, both leading to more jets in an event. Therefore, the analysis targets signatures with a range of jet multiplicities, ranging from at least 2 to at least 6 high p_T jets. Note that although these SUSY models are specifically targeted, the analysis is sensitive to any model with strongly-interacting particles decaying to two or more jets and undetectable particles.

The Standard Model backgrounds which have the signature of high p_T jets and missing transverse momentum are QCD multi-jets, the production of top quarks (either alone or in $t\bar{t}$ pairs), vector boson production with additional jets, and dibosons. These will be described in more detail in section 5.1.1. By targeting a fully hadronic final state, without any leptons, backgrounds which have real $\cancel{E}_T$ together with a lepton, such as W+jets with leptonic W decay, are suppressed. Since hadronically decaying tau leptons are difficult to identify, they are not vetoed. A selection without light leptons does come at the cost of high background yields for other processes: QCD multi-jet events can become a background, even though these events have low $\cancel{E}_T$, due to the large QCD cross section at the LHC. The distribution of $\cancel{E}_T$ and p_T of the leading jet in events with at least 2 jets is shown in figure 5.1. The distributions for two squark-pair production models are shown in the dashed histograms for a model with a small mass splitting between the squark and LSP (in light blue) and with higher squark mass and lower LSP mass (in red). Both $\cancel{E}_T$ and the leading jet p_T are

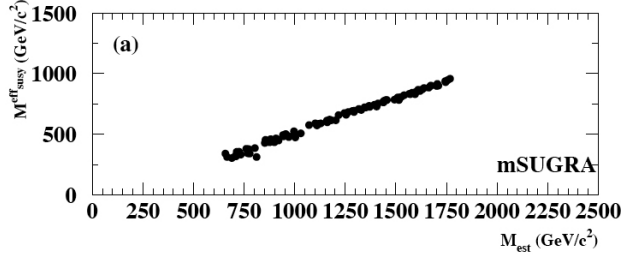


Figure 5.3: The correlation between the estimated SUSY mass scale M_{est} , obtained from the mean of the m_{eff} distribution from equation 5.1, and the effective SUSY mass scale M_{SUSY}^{eff} in mSUGRA models. Figure taken from [229].

correlated to the mass splitting of the SUSY scenario, with on average lower values for low mass-splitting scenarios, or ‘compressed’ spectra. The distribution of $\cancel{E}_T$ and jet p_T of these compressed SUSY models look much like those of the SM backgrounds in the case of direct decay of the heavy SUSY particles to jets and LSPs. This is demonstrated as well in figure 5.2, which shows an illustration of the typical $\cancel{E}_T$ and leading jet p_T for events in simplified direct decay SUSY models with varying levels of mass splittings, and of the SM backgrounds. The same holds for SUSY models with gluino production. As compressed scenarios look so much alike to SM backgrounds, they are difficult to discover or exclude. A dedicated search is needed to be sensitive to them, which will be presented later in this chapter. The simplified models, which do not have any constraints on their mass spectra, are used to interpret the results for both high mass scenarios, and compressed spectra. The CMSSM on the other hand has by definition a large mass splitting between the squarks and gluinos on the one hand, and the LSP on the other. It is therefore used to study high mass scenarios with possibilities of long decay chains.

To discriminate signal and background events, the effective mass m_{eff} variable is used [230], which is defined as the scalar sum of the missing transverse energy and the transverse momenta of the jets in the event:

$$m_{eff} = \cancel{E}_T + \sum_{j=1}^n p_T^j, \quad (5.1)$$

with the sum running over all n jets in the event. It has been shown [229] that when no other objects are present in an event apart from the selected jets, the effective mass is strongly correlated with the mass of the produced SUSY particle-pair, which is illustrated in figure 5.3; in this figure, M_{est} represents our definition of m_{eff} , while M_{SUSY}^{eff} is an effective SUSY mass scale, taking the production cross sections of the produced sparticles into account:

$$M_{SUSY}^{eff} = M_{SUSY} - \frac{m_{\tilde{\chi}_1^0}^2}{M_{SUSY}}, \quad (5.2)$$

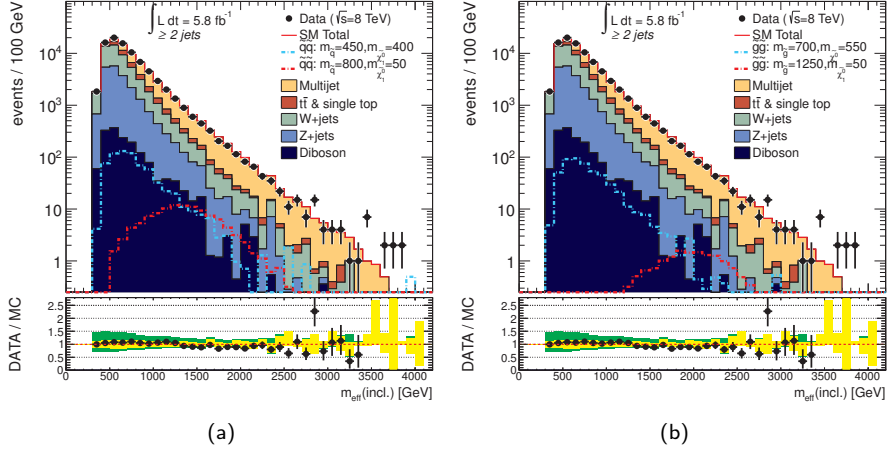


Figure 5.4: Observed and simulated m_{eff} distribution in events with no electrons or muons, $\cancel{E}_T > 160$ GeV and at least 2 jets (with $p_T > 130$ and 60 GeV, respectively) in 5.8 fb^{-1} of $\sqrt{s} = 8$ TeV data. The dashed histograms show the distribution of two signal models: (a) squark-pair production decaying directly, with $m_{\tilde{q}} = 450$ GeV, $m_{\tilde{\chi}_1^0} = 400$ GeV (light blue) and $m_{\tilde{q}} = 800$ GeV, $m_{\tilde{\chi}_1^0} = 50$ GeV (red), and (b) gluino-pair production decaying directly, with $m_{\tilde{g}} = 700$ GeV, $m_{\tilde{\chi}_1^0} = 550$ GeV (light blue) and $m_{\tilde{g}} = 1250$ GeV, $m_{\tilde{\chi}_1^0} = 50$ GeV (red).

with M_{SUSY} defined as the cross section weighted SUSY mass scale:

$$M_{\text{SUSY}} = \frac{\sum_i \sigma_i m_i}{\sum_i \sigma_i}. \quad (5.3)$$

Here the σ_i and m_i are the cross section and average mass of the pair of initially produced particles, respectively. Because SUSY particles are expected to have high mass, m_{eff} can be a powerful discriminant between signal and SM backgrounds.

A m_{eff} distribution is shown in figure 5.4, which compares data to the MC background predictions. Two signal distributions for two squark-pair production models with a different $\Delta m(\tilde{q}, \tilde{\chi}_1^0)$ are shown in the left figure, while the right figure shows the same data and backgrounds, yet with distributions for two gluino-pair production models. Although m_{eff} is correlated to the squark or gluino mass, it decreases in models with high LSP mass.

The analysis presented in this thesis is set up as a counting experiment: upon each event a set of requirements (cuts) is placed. These requirements are optimised on simulated data to reject SM backgrounds, while accepting the SUSY signal as efficiently as possible. Events which pass all requirements are said to be signal-like, and the final set of cuts are called the *Signal region* (SR) selection. Analyses which

focus on more than one signal model or more than one signature need separate signal regions for each of these signatures. In our case, there will be 5 SR channels for each of the 2-6 jet multiplicities.

To estimate the number of SM background events in the SRs, special control regions (CRs) are defined for each major background, enriched in a particular background process. The set of cuts which determine these CRs are defined as closely to the SR selections as possible, while adding requirements to select the background process in question. This means that each individual background has 5 CRs corresponding to the 5 jet multiplicity channels. The CRs are used as a semi-data driven estimation of the background in the SRs: the SM background estimation is obtained from simulation, but is subsequently normalised using the observed data in the CRs. The normalisation is performed using transfer factors (TFs) between the SR and CRs. For each background process p the transfer factor from its CR to the SR is defined as being the ratio of the expected number of events in these regions:

$$TF_p = \frac{N_{SR_p}^{exp}}{N_{CR_p}^{exp}}. \quad (5.4)$$

Here $N_{SR_p}^{exp}$ and $N_{CR_p}^{exp}$ are the number of expected events for process p from simulation in the SR and CR, respectively. Using these transfer factors and the number of observed events in the CRs, a semi-data driven estimate of the number of events in the SR for this process is obtained: $N_{SR_p}^{predicted} = TF_p \cdot N_{CR_p}^{observed}$. Furthermore, transfer factors between the different CRs allow for a coherent normalisation across all regions. The transfer factors are very useful in the uncertainty calculations, as many systematic uncertainties which are correlated between SR and CRs, such as the jet energy scale (JES), largely cancel in the ratio 5.4. The background estimation is described for each background separately in section 5.4.

After counting the number of observed events in the SR and comparing to the estimated number of SM background events, either a discovery can be claimed, or an upper limit on the possible number of SUSY events and cross section of allowed SUSY production can be placed using the statistical procedure. This upper limit is interpreted as an exclusion limit in a SUSY mass plane.

5.1.1 Standard Model backgrounds to hadronically decaying SUSY

As mentioned in the introduction of this chapter, many Standard Model processes will act as a background in the search for SUSY in hadronic signatures. As the selected signature of these SUSY events is a number of high momentum jets and missing transverse momentum, while events with leptons (electrons and muons) or photons are discarded, Standard Model backgrounds to this SUSY signature will thus be processes with real or fake $\cancel{E}_T$ and high p_T jets, which either produce only jets, or also photons or tau leptons, or produce electrons or muons which are not identified by the detector. Hadronically decaying tau leptons are difficult to identify and are observed as jets, and they are therefore not vetoed. The most important backgrounds to this analysis are discussed below.

QCD multi-jets

The most abundant process produced in pp collisions in the LHC is QCD multi-jet production with a cross section of $\mathcal{O}(10^{10})$ pb. QCD multi-jets are the result of parton scattering: two-parton scattering can result in a high p_T di-jet event, while hard gluon emission can lead to several additional high p_T jets. Although this process does not generate 'real' missing transverse momentum, the misreconstruction of jets can lead to a net $\cancel{E}_T$. This form of $\cancel{E}_T$ will thus usually be aligned with one of the hardest jets.

Although the cross section is smaller, parton scattering can also occur with heavy flavour quarks. Another source of $\cancel{E}_T$ can then be semi-leptonic decay of these heavy flavour quarks, leading to a lepton and a neutrino. If the electron or muon is either soft, outside of the detector acceptance or if the lepton is reconstructed as a jet, this leads again to a signature of high p_T jets with $\cancel{E}_T$.

W+jets

When W bosons are produced in collisions, they decay either into hadrons, $W^\pm \rightarrow q\bar{q}'$, with a branching ratio of 67.6%, or into leptons, $W^\pm \rightarrow l^\pm \nu_l$, with $l = e, \mu, \tau$, with a branching ratio of 10.8% per flavour [19]. Such a production of a W boson can be accompanied by additional jets, where the production cross section decreases roughly by $1/\alpha_S$. In the case of a decay to a hadronically decaying τ , or with a misidentified e or μ , a SUSY signature is obtained, as the boosted neutrino again leads to real $\cancel{E}_T$ in the detector.

Z+jets

Just like W boson production, Z boson production is an important background for SUSY searches when accompanied by additional jets. Z bosons decay either hadronically (BR = 70%), into two leptons, $Z \rightarrow l^+ l^-$ (BR = 10% for the sum of all lepton flavours), or into two neutrinos, $Z \rightarrow \nu\nu$ (BR = 20%). The latter decay is an irreducible background to our search, as it has the exact same signature when accompanied by hard jets.

Diboson

The production of two vector bosons, WW , WZ or ZZ , shows similar behaviour to W and Z plus jets. Diboson production has a much lower cross section than the production of a single W or Z , yet at the luminosity achieved in 2012 it is not a negligible background.

Top quark backgrounds

Top quarks can be produced either in pairs ($t\bar{t}$) or on their own. Top pair production occurs either via quark-antiquark interaction ($q\bar{q} \rightarrow t\bar{t}$) or via gluon fusion ($gg \rightarrow t\bar{t}$).

Since anti-quarks are only available via sea quarks, gluon fusion is the dominant production mechanism at the LHC. Top quarks decay before they can hadronise, into a W boson plus a b quark ($|V_{tb}| > 0.999$ [19]), with a minimal number of top quarks decaying into a W boson with a light-flavour quark (d or s). Due to the W decay modes discussed above, $t\bar{t}$ pairs can decay either fully hadronically, semileptonically when only one W boson decays into leptons, or dileptonically. Semileptonic $t\bar{t}$ decay thus can lead to real $\cancel{E}_T$, at least four jets and a lepton. ISR and FSR can lead to additional jets in the event. If the lepton is either a τ or is misidentified, it leads to the signature under study.

Single top production occurs either together with a W boson or with another quark via s - and t -channel diagrams. Again, the top quark decays into a W boson plus a b quark.

All other backgrounds are negligible due to either the event selection or a too small cross section. For each of these backgrounds except the diboson background a separate CR has been defined, as will be described further in section 5.4.

5.2 Dataset and Monte Carlo samples

The data samples used in this analysis were gathered by the ATLAS detector in early 2012 and add up to an integrated luminosity of 6.0 fb^{-1} . After the application of data quality requirements, via a so-called Good Run List (GRL), 5.8 fb^{-1} remains. Using a technique similar to that of refs. [231, 232], the uncertainty on the integrated luminosity is measured to be 3.6%. The GRL lists luminosity blocks without detector malfunctions or inefficiencies in any of the sub-detectors. When vetoing luminosity blocks not in the GRL, impurities in the used data are reduced significantly.

5.2.1 Monte Carlo Samples

To compare the recorded data to theoretical predictions, Monte Carlo simulations are used. The simulation steps and various available generators have been described in section 4.7. This section gives the details on the used generators for each of the background and signal MC samples. The MC generators of the backgrounds are summarised in table 5.1.

Simulated W +jets events are generated in this analysis using SHERPA 1.4.0 [43], with up to five additional partons generated in the matrix element. The MENLOPS [233, 234] technique is used to match the matrix element jets to those of the parton showers by the CKKW scheme [210–212]. The SHERPA samples use the NLO CT10 PDFs, while the generator is able to perform fragmentation and parton showering by itself. Both b - and c -quarks are generated as massive particles. For $W^\pm (\rightarrow l^\pm \nu)$ +jets production, 40 million events are generated for each lepton flavour. As only a very small number of these survive the SUSY SR selection, leading to low MC statistics, additional events were generated with a higher probability to survive the SR cuts: 1 million events were generated per lepton flavour requiring the W boson to have at least $p_T^W > 200 \text{ GeV}$ at truth level, while 4 million events were generated per

Sample	MC generator	Comments
W +jets	SHERPA	Alternative: ALPGEN
Z/γ^* +jets	ALPGEN	Alternative for $Z(\rightarrow l^+l^-)$ +jets: SHERPA
γ +jets	ALPGEN	-
$t\bar{t}$	MC@NLO	Alternative: SHERPA
single top: s - and Wt -channel	MC@NLO	
single top: t -channel	ACERMC	Alternative: SHERPA
Diboson	SHERPA	-
QCD	PYTHIA	-
CMSSM	HERWIG++	$\tan\beta = 10$, $A_0 = 0$, $\mu > 0$ $m_0 \in [200, 4000]$ GeV, 200 GeV steps, $m_{1/2} \in [250, 900]$ GeV, 50 GeV steps
Squark-pair production, and direct decay	MADGRAPH+PYTHIA	$m_{\tilde{q}} \in [87, 1575]$ GeV $m_{\tilde{\chi}_1^0} \in [0, 1200]$ GeV
Gluino-pair production, and direct decay	MADGRAPH+PYTHIA	$m_{\tilde{g}} \in [87, 1800]$ GeV $m_{\tilde{\chi}_1^0} \in [0, 1200]$ GeV
Squark-gluino production, and direct decay	MADGRAPH+PYTHIA	$m_{\tilde{q}} = 0.96 \times m_{\tilde{g}} \in [87, 1967]$ GeV $m_{\tilde{\chi}_1^0} \in [0, 1200]$ GeV

Table 5.1: Monte Carlo generators for the SM backgrounds and SUSY signal samples used in this analysis. The alternative samples quoted are used for the evaluation of theoretical uncertainties. For the simplified models, the step size is variable and thus not specified.

flavour requiring at least 3 truth jets in the final state. The former samples enhance the statistics in the low jet multiplicity SRs, while the latter samples will increase the statistics in higher jet multiplicity samples. At the analysis level all samples are combined using truth level information.

For Z/γ^* +jets and γ +jets samples the ALPGEN 2.14 [217] generator is used², with again up to five additional partons generated in the matrix element. Parton showering is simulated by HERWIG [213, 214] with JIMMY [216] modelling the underlying event, while the CTEQ6L1 PDF set [235] is used. The matching of the matrix element partons to the parton shower jets (see section 4.7.1) is done using the MLM matching scheme [208]. For $Z(\rightarrow \nu\nu)$ +jets and γ +jets the samples are sliced in bins of vector boson p_T^V , such as to have enough statistics in the tails of the momentum distributions. As the momentum cut increases from $p_T^V > 35$ GeV to $p_T^V > 500$ GeV, the equivalent integrated luminosity ranges from 1 fb^{-1} to 500 fb^{-1} . $W^\pm(\rightarrow l^\pm\nu)$ +jets, $Z(\rightarrow$

² Z/γ^* +jets denotes the production of Z bosons, as well as Drell-Yan processes, where an off-shell photon acts as propagator.

l^+l^-)+jets with $l = e, \mu$, and $Z(\rightarrow \nu\nu)$ +jets are normalised to the NLO predictions of 11.88 fb, 1.21 fb and 5.68 fb respectively, while the sliced γ +jets samples are each normalised to the cross sections obtained from ALPGEN.

The main theoretical uncertainties on the W +jets, Z/γ^* +jets and γ +jets samples affect the shape of the kinematic distributions of variables used in the SUSY search (since the MC is normalised to data via the CRs), and are caused by: dependence on the choice of the renormalisation and factorisation scale; the choice of the energy scale of the matching between matrix element and parton shower; the number of additional partons which are generated; and uncertainties from the PDFs. These are all covered in section 5.4.6.

$t\bar{t}$ samples are generated with MC@NLO [218, 219], interfaced with HERWIG simulating the parton showers and JIMMY to model the underlying event. The PDFs which are used are from the CT10 PDF set. The $t\bar{t}$ cross section is predicted to be 238^{+22}_{-24} pb from approximate NNLO QCD calculations [236]. The main theoretical uncertainty on this sample is due to the number of additional jets, as MC@NLO only generates one additional parton in the matrix element. This uncertainty is evaluated by comparing to a SHERPA $t\bar{t}$ sample.

Both the s - and Wt -channel of the single top MC samples are generated in the same manner as $t\bar{t}$ discussed above. For the t -channel ACERMC is used, interfaced with PYTHIA 6 [142], due to mismodelling in MC@NLO when a b -quark exists in the initial state. The cross sections for these processes are 5.61 ± 0.22 pb, 22.4 ± 1.5 pb and $87.8^{3.4}_{-1.9}$ pb for s -channel, Wt -channel and t -channel respectively.

SHERPA is used to generate WW , WZ , ZZ , $W\gamma$ and $Z\gamma$ diboson processes, while QCD multi-jet MC samples are generated with PYTHIA8 [237], using CT10 PDFs and the AU2 tune [238] by ATLAS. As a fully data-driven method is implemented to estimate the QCD multi-jet background, these latter samples are only used for optimisation studies and comparison between data and MC.

The choice for the above generators is made after comparison to data in the control regions. ALPGEN describes the Z/γ^* +jets data better, while the available ALPGEN W +jets sample has too low statistics leading to the choice for SHERPA for W +jets. For Z/γ^* +jets, W +jets and $t\bar{t}$ simulations alternative samples are defined as well for the evaluation of theoretical uncertainties, as is discussed in section 5.4.3.

SUSY signal MC

Since the analysis targets the CMSSM and certain simplified models, see section 1.3.4, with the focus on the more compressed regions, events need to be simulated for these SUSY signals as well. For each of the models, a set of points is defined in a 2- or 3-dimensional parameter phase space. In the CMSSM this is m_0 versus $m_{1/2}$, while the SUSY particle masses are used for the simplified models. For each of the points in these grids, between 5000 and 20000 events are generated, depending on the expected signal acceptance.

For the CMSSM, the mass spectrum and decay modes of the SUSY particles are

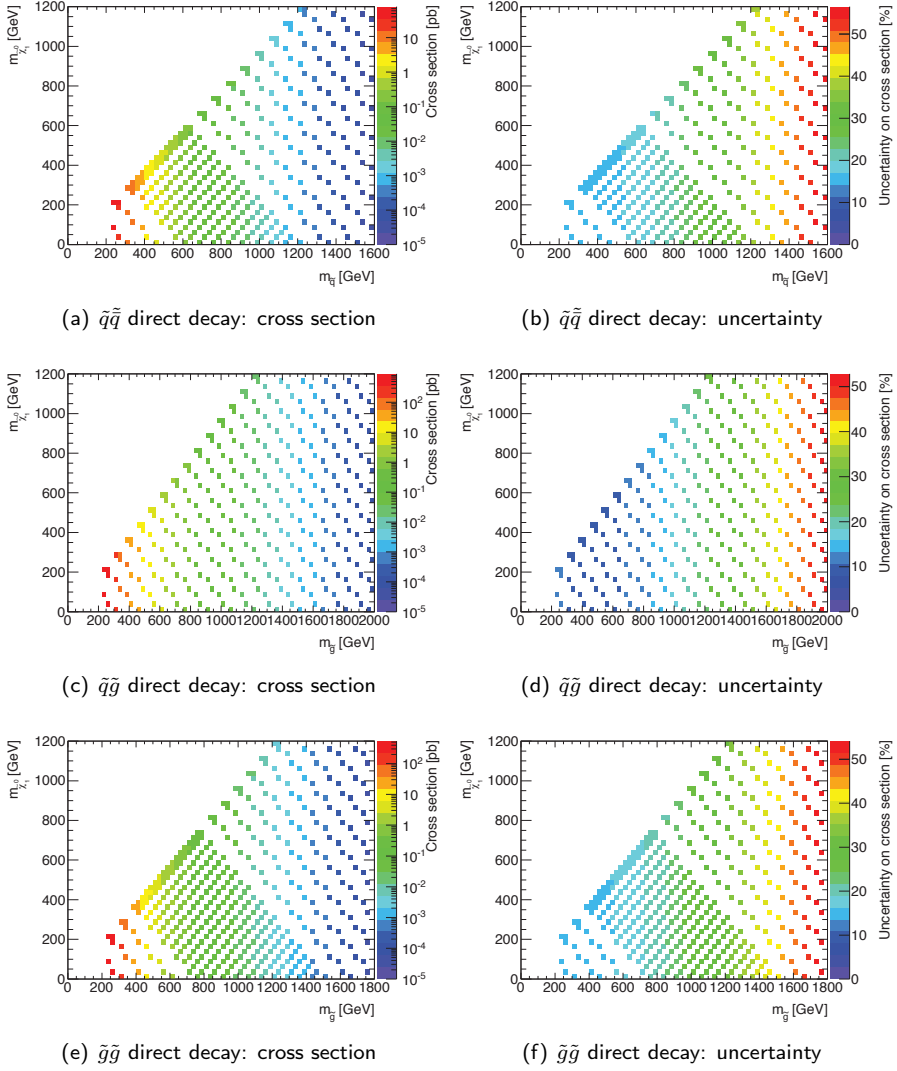


Figure 5.5: Cross section in pb (left) and its theoretical uncertainty in % (right) for: on the top row the $q\bar{q}$ production grid; on the middle row the $q\bar{q}$ production grid; and on the bottom row the $g\bar{g}$ production grid.

generated using SUSY-HIT [239], while the events are generated using HERWIG++ 2.5.2 [215] with the UE-EE-3 tune [240]. All simplified model MC samples are generated using MADGRAPH5 [226] with one additional parton in the matrix element, while the parton showering is performed by PYTHIA 6. The MLM matching scale is set to 50 GeV. This value is varied by a factor of 2 to evaluate the systematic uncertainty

related to ISR jets, as described in section 5.4.7.

The choice for MADGRAPH stems from the fact that it is a matrix element generator, which is able to add one or more extra outgoing partons to the matrix element, corresponding to possible ISR or FSR jets. In the case of HERWIG++, extra partons are added using parton showering, which does not describe hard ISR or FSR radiation well. Furthermore, the uncertainty estimation due to ISR production is easier evaluated via MADGRAPH. As we will see, the ISR production is important for compressed spectra. For both CMSSM and simplified models the used PDFs are from the CTEQ6L1 set. Unlike the CMSSM signal samples, the detector simulation for simplified model samples is not performed by a full simulation, but by a fast ATLAS simulation using ATLFast II [241]. The cross sections are obtained for each process and signal grid individually as described in chapter 3. The cross sections and their theoretical uncertainties for the direct decay simplified models are shown in figure 5.5. For the CMSSM, the cross sections and uncertainties on the squark and gluino production processes can be found in figures 3.9 and 3.10.

5.3 Event selection of the 0-lepton analysis

After the motivation and brief overview of the analysis given in the introduction, and the description of the recorded data and used MC samples in the previous section, all the prerequisites are met to be able to discuss the analysis in detail. Both data and simulated events need to pass an event selection, with requirements on the event properties to categorise it as being signal-like (it falls in the SR) or background-like (it falls in a CR).

5.3.1 Trigger and event selection

The trigger used for the event selection in the signal regions is a combined jet plus $\cancel{E}_T$ trigger, EF_j80_a4tchad_xe100_tclcw_veryloose. It selects events with at least one jet with $p_T > 80$ GeV, which is reconstructed from topological clusters using an anti- k_T algorithm with distance parameter $R = 0.4$, together with missing transverse momentum of at least 100 GeV. When applying cuts of $p_T > 130$ GeV and $\cancel{E}_T > 160$ GeV in this analysis, the trigger is 100% efficient as is shown in figure 5.6.

For the control regions, different triggers are applied to be able to select events with leptons and photons. In the control regions of W +jets and $t\bar{t}$ backgrounds, events are triggered using the trigger chains with the lowest available single lepton p_T triggers: for muons the EF_mu24i_tight trigger, while electrons are triggered with EF_e24vhi_medium1 and EF_e60_medium1. Photons are triggered using EF_g120_loose in the control region targeting photons.

Event cleaning

After passing the trigger, a set of ‘cleaning cuts’ are placed on the events to remove sources of possible fake missing transverse momentum. Such sources are non-

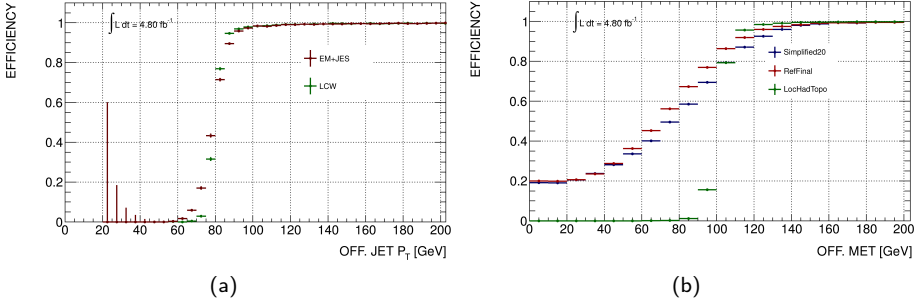


Figure 5.6: Efficiency of the *EF_j80_a4tchad_xe100_tclcw_veryloose* trigger for (a) the leading jet p_T for the LCW and EM+JES jet calibrations, and (b) three E_T definitions, in a data sample selected with a muon trigger. In this thesis the LCW jet calibration and RefFinal E_T definition are used, and for the E_T term a leading jet cut of $p_T > 130$ GeV is applied. Figure taken from [242].

functioning calorimeter cells, misidentified muons and non-collision backgrounds, i.e. particles in the detector from other sources than the pp collision, for instance beam-gas events. These are all shortly discussed in the following paragraphs.

To remove events contaminated by noise in the liquid argon calorimeter, events which are identified to have experienced significant noise in the LAr EM calorimeter ($\text{larerror} \neq 0$) are vetoed. Jets coming from non-collision backgrounds, such as cosmic rays, beam halo events or calorimeter noise, are rejected using jet quality requirements [243], discussed in the following. Calorimeter noise can lead to fake energy deposits, which can be reconstructed as jets. However, deposits from real particle showers cause a characteristic pulse in the calorimeter cells, which can be distinguished from noise. For each pulse the quadratic difference $Q_{\text{cell}}^{\text{LAr}}$ between the measured and expected pulse is used to discriminate noise from real particle showers. On jet level, the normalised average jet quality $\langle Q \rangle$ and fraction of energy in the LAr and HEC calorimeters with low signal shape quality ($Q_{\text{cell}}^{\text{LAr}} > 4000$, f_Q^{LAr} and f_Q^{HEC} respectively, are used. To remove sporadic noise bursts in the hadronic end-caps, we require that jets are removed from our selection if either ($f_{\text{HEC}} > 0.5$ and $|f_Q^{\text{HEC}}| > 0.5$ and $\langle Q \rangle > 0.8$) with f_{HEC} the energy fraction deposited in the HEC, or a large negative energy $|E_{\text{neg}}| > 60$ GeV, since the neighbouring cells will appear to have negative energy due to capacitive coupling between the cells. Coherent noise in the LAr calorimeter is reduced by vetoing jets with $f_{\text{em}} > 0.95$ and $f_Q^{\text{LAr}} > 0.8$ and $\langle Q \rangle > 0.8$, with again f_{em} being the energy fraction deposited in the electromagnetic calorimeter. The efficiency of these requirements is 99.8%, while the fake jet rejection is high.

Non-collision backgrounds which do not have a track are suppressed by requiring that the charged fraction ($f_{\text{ch}} = \sum |p_T^{\text{trk}}|/p_T^{\text{jet}}$) of the leading two jets is not too small. If the electromagnetic contribution to the jets, f_{em} , is significant, the cut on f_{ch} may

be tightened since it is likely that there are fewer neutrons in the event. The final requirements are therefore: jets in $|\eta| < 2.0$ are rejected if $f_{ch} < 0.02$ or if both $f_{ch} < 0.05$ and $f_{em} > 0.9$. The cut is found to be very efficient, with more than 99% of events being accepted while rejecting cosmic events, detector noise and beam backgrounds.

The Tile calorimeter very occasionally experiences data corruption, which can lead to large negative energy in cells and clusters, and thus large fake $\cancel{E}_T$. Since these noisy cells, by construction, do not contribute to any object (e.g. a jet), they are added to the $\cancel{E}_T^{CellOut}$ term. A requirement is placed such that the $\cancel{E}_T^{CellOut}$ term cannot contribute too much to the total $\cancel{E}_T$; events are rejected if

$$\frac{\cancel{E}_T^{CellOut}}{\cancel{E}_T} \cos\left(\phi(\cancel{E}_T^{CellOut}) - \phi(\cancel{E}_T)\right) > 0.5. \quad (5.5)$$

Here the cosine over the angle ensures we compare only the $\cancel{E}_T^{CellOut}$ term in the direction of the total $\cancel{E}_T$. The inefficiency of the cut is found to be negligible. Two cells of the Tile calorimeter have been found to be malfunctioning, while not being removed in reconstruction. Events with a jet pointing at such a cell are rejected if a large fraction of energy is deposited in the malfunctioning layer of the cell, $E_{layer}/E_{jet} > 0.6$, or if any of the two leading jets pointing to the region of the affected cells has $f_{ch} < 0.2$ and $f_{em} < 0.3$.

‘Fake’ muons, i.e jets mistakenly reconstructed as muons, can have a large contribution to $\cancel{E}_T$. They come from jets which punch-through the calorimeter to spill over in the muon spectrometer or from inner detector tracks which are not correctly reconstructed and matched to tracks in the muon spectrometer. Moreover, muons which are badly reconstructed with a reconstructed momentum far away from the true value (usually due to momentum mismeasurements in the MS) could give high fake $\cancel{E}_T$. To solve this, events are rejected if they have a muon whose error on the momentum is large with respect to its momentum: $\sigma(q/p)/|q/p| > 0.2$, where q/p is the ratio of charge over momentum, which is measured by the MS for a muon travelling through the magnetic field. Moreover, to have a higher rejection efficiency, a cut is placed on the muon contribution $(\cancel{E}_T)^{Muon}$ to the total $\cancel{E}_T$ in the same form of equation 5.5. Events with cosmic muons, which could contaminate the CRs, are rejected by requiring the muons to come from the primary vertex: muons are required to have $|d_0| \leq 0.2$ mm and $|z_0| \leq 1$ mm, with d_0 and z_0 the transverse and longitudinal distance to the primary vertex, respectively.

Finally, fake missing momentum can also be caused by jets falling in non-operational calorimeter cells. The Tile calorimeter had between 0.5 and 1% dead cells during the data taking periods used in this analysis. Events with jets falling into a region with a dead Tile cell are rejected if a jet is close to the $\cancel{E}_T$ direction in ϕ and the fraction of energy it loses due to the malfunctioning cell is large: $\Delta\phi(j, \cancel{E}_T) < 0.2$ for jets with $p_T > 40$ GeV and $B_{corr}^{jet} > 0.05$. Here B_{corr}^{jet} is the estimated fraction of lost energy, estimated by profiling the jet shape.

The inefficiency of the last requirement is $< 2\%$ for SUSY signal, while for all

other cleaning cuts the inefficiency is proven to be less than one percent; see for instance [243].

Event preselection

Apart from the trigger and cleaning cuts, events are required to have no electrons or muons. This requirement rejects events coming from W and $t\bar{t}$ production decaying leptonically. To achieve this, events with electrons with $p_T > 20$ GeV are vetoed, while, due to a better understood veto efficiency, events with muons are rejected from $p_T > 10$ GeV onwards. Note that hadronically decaying τ -leptons are not vetoed. A veto on events with isolated tracks has been studied to remove tau candidates, which seems to give a minor gain in sensitivity. This veto might be implemented in the analysis on a larger dataset.

To ensure the events are on the ‘trigger plateau’, where the trigger efficiency is approximately 100%, the fully reconstructed leading jet p_T and $\cancel{E}_T$ are required to be higher than 130 GeV and 160 GeV respectively. This is seen from the efficiency turn-on curves shown in figure 5.6.

As mentioned in the first section, signal events are expected to have at least 2–6 jets, depending on the produced SUSY particles. Therefore five analysis channels are defined, characterised by the jet multiplicity: channel A, B, C, D and E contain events with 2 or more, 3 or more, 4 or more, 5 or more and 6 or more jets, respectively. Many SM backgrounds are reduced when requiring more jets, as each additional jet reduces the corresponding cross section by approximately a factor of $1/\alpha_S$. To reject contamination of pile-up jets, which have predominantly low p_T , each of the leading N jets in the inclusive N -jet channel is required to have $p_T > 60$ GeV, while the additional jets should at least obey $p_T > 40$ GeV.

QCD multi-jet events are a large background to our SUSY signals. When a jet is mismeasured, ‘fake’ $\cancel{E}_T$ is generated, typically pointing in the direction of the mismeasured jet. Furthermore, the production of a heavy-flavour quark inside a jet with subsequent semileptonic decay will also lead to $\cancel{E}_T$ pointing along the jet. Therefore, in order to reject these events, a cut is placed on the minimum ϕ angle between the jets and the $\cancel{E}_T$. Figures 5.7 (a) and (b) show the distribution of $\min(\Delta\phi(j_i, \cancel{E}_T))$ for the three jets with highest p_T in the inclusive 2-jet channel and the 5-jet channel, with a large peak at small values coming from QCD multi-jet events. Requiring $\Delta\phi(j_i, \cancel{E}_T) > 0.4$ for the leading three jets i , and $\Delta\phi(j_i, \cancel{E}_T) > 0.2$ for the remaining jets is seen to suppress the QCD background with a minimal loss of signal events.

To further suppress backgrounds without a source of ‘real’ $\cancel{E}_T$ such as QCD multi-jets, a requirement on the $\cancel{E}_T$ fraction of the effective mass is placed. As these type of events usually have less $\cancel{E}_T$, it is expected that $\cancel{E}_T$ contributes less to the total m_{eff} than high m_{eff} events with a real source of $\cancel{E}_T$. However, instead of cutting on $\cancel{E}_T/m_{\text{eff}}$, a slight improvement is made. Decays from SUSY particles often produce more jets than the SM background events, increasing the total m_{eff} , and therefore decreasing the ratio. If we however use the *exclusive* effective mass $m_{\text{eff}}^{\text{excl}}$ in the ratio,

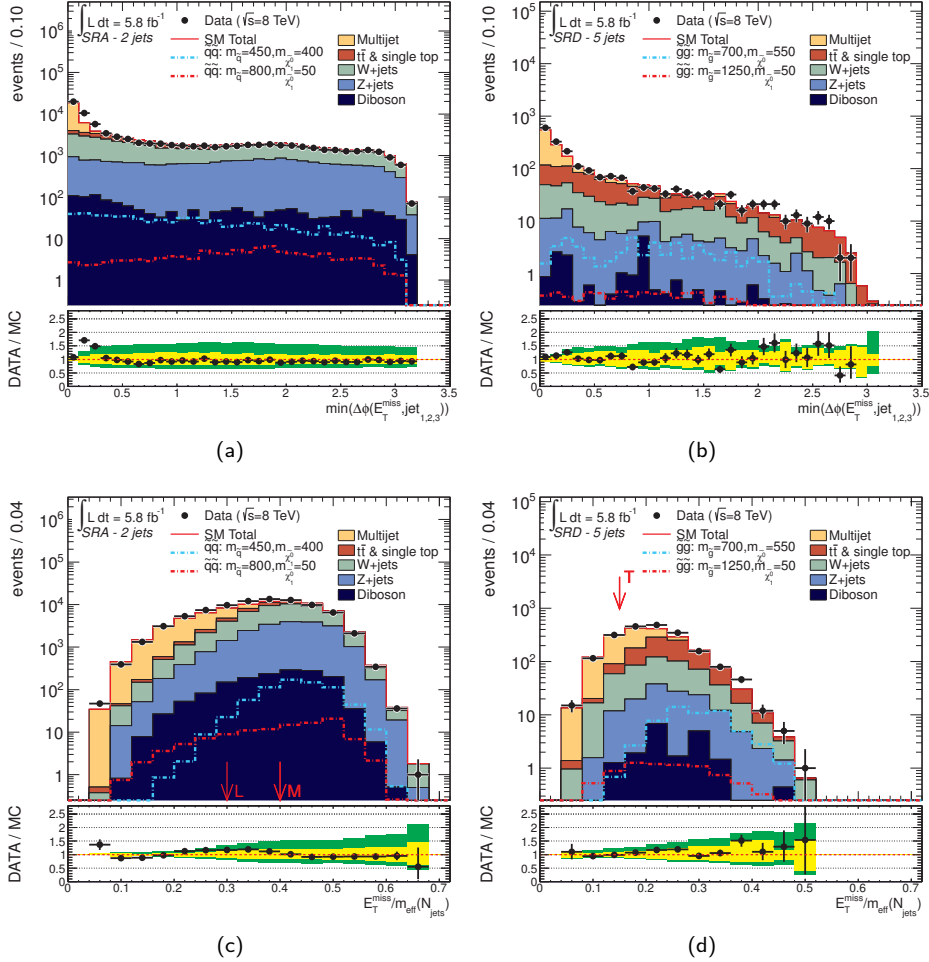


Figure 5.7: Top row: observed distributions of the minimum of the ϕ angle between the three leading jets and $\cancel{E}_T$. Bottom row: observed $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ distributions. Events are selected with no electrons or muons, $\cancel{E}_T > 160$ GeV, while the $\Delta\phi(j_i, \cancel{E}_T)$ cut has not yet been applied. Left figures (a and c) show events with at least 2 jets, right figures (b and d) with at least 5 jets. Data is compared to backgrounds from MC simulations and pseudodata in case of QCD multijets. The signal models shown as the dashed histograms are from the squark-pair model in the 2-jet selection, while the 5-jet selection has two gluino-pair signal distributions. The red arrows indicate the SR selections, where the “L”, “M” and “T” stand for loose, medium and tight selection.

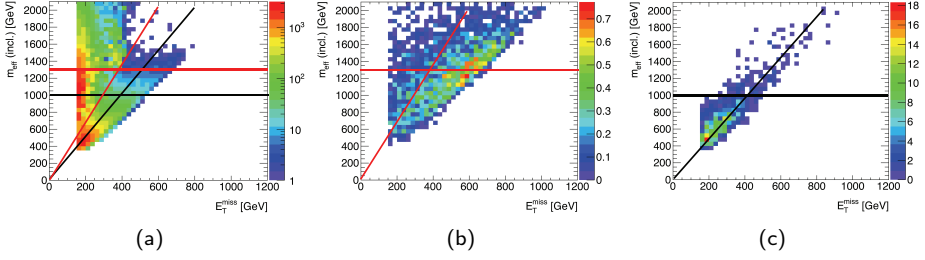


Figure 5.8: Two dimensional distributions of $\cancel{E}_T$ and m_{eff} in simulated events with at least 2 jets after the $\Delta\phi$ cut, for (a) the sum of all background events, (b) squark-pair production with $m_{\tilde{q}} = 800$ GeV and $m_{\tilde{\chi}_1^0} = 50$ GeV and (c) with $m_{\tilde{q}} = 450$ GeV and $m_{\tilde{\chi}_1^0} = 400$ GeV. The red lines show the cuts used in the ‘tight’ selection, while the black lines represent the ‘loose’ selection. Note that the z -axis is on a logarithmic scale for the backgrounds in (a).

defined in an N -jet channel as $m_{\text{eff}}^{\text{excl}} = \cancel{E}_T + \sum_{j=1}^N p_T^j$ (in contrast to the standard definition of m_{eff} which takes into account all jets), the signal efficiency of such a requirement increases with respect to the background efficiency [244]. In other words, a cut on $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ is still effective against fake sources of $\cancel{E}_T$, while more SUSY signal events will pass the requirement.

Figures 5.7 (c) and (d) show the distribution of $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ for the 2- and 5-jet channel. In the 2-jet distribution, two signal samples have been overlaid from a squark-pair production model with a small mass splitting in light blue and a large mass splitting in red. For the 5-jet channel two gluino-pair signal samples are shown. The $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ distribution is shifted to lower values for these events with higher jet multiplicities. This can be explained as follows: within SUSY, the possibility of a low p_T LSP increases in events with a higher number of jets, decreasing $\cancel{E}_T$ in these events. On the other hand, $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ is slightly higher for compressed SUSY than for larger mass splittings. Compressed SUSY signals have lower $\cancel{E}_T$ as well as softer jets and thus m_{eff} , which leads to a higher value of $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$. Since the QCD cross section decreases as well with increasing jet multiplicity, we expect that the $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ requirement may be loosened for the higher jet multiplicity channels.

Finally, a requirement on the effective mass is placed to discriminate signal events from the remaining backgrounds. As explained earlier, m_{eff} peaks around the sum of the masses of the produced sparticles, thus the cut on m_{eff} can be harder for SRs targeting scenarios with higher mass SUSY particles. The optimisation of the final values of the $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ and m_{eff} cuts is discussed in the following section for compressed SUSY models – the optimisation for high mass models with low compression is performed similarly.

In figure 5.8 both $\cancel{E}_T$ and m_{eff} are shown for simulated events with at least 2 jets and passing the $\Delta\phi$ requirement. The required values for m_{eff} and $\cancel{E}_T/m_{\text{eff}}$ are shown

Cut	Description	A'	A	B	C	D	E
			(l/t)		(l/m/t)		(l/m/t)
Lepton veto		No events with e or μ after overlap removal					
$\cancel{E}_T$ [GeV] >					160		
Jet 1 p_T [GeV] >					130		
Jet multiplicity		≥ 2	≥ 2	≥ 3	≥ 4	≥ 5	≥ 6
Jet p_T [GeV] >			60 (for jets 2-4)			40 (for jets 5 and 6)	
$\Delta\phi(j_i, \cancel{E}_T)$ $i \leq 3$			0.4 for jets with $p_T > 40$ GeV				
$\Delta\phi(j_i, \cancel{E}_T)$ $i > 3$		–	–	–	0.2	0.2	0.2
$\cancel{E}_T/m_{\text{eff}}^{\text{excl}} >$		0.4	0.3	0.25	0.25	0.2	0.15
m_{eff} [GeV] >		1200	1400/1900	1900	900/1200/1500	1500	900/1200/1400

Table 5.2: Event selection used in the $\sqrt{s} = 7$ TeV 2011 analysis from the lepton veto onwards, with the separation between ‘loose’ (l), ‘medium’ (m) and ‘tight’ (t) in the last m_{eff} cut.

for the ‘tight’ selection in red lines, and in black for the ‘loose’ selection. Note that the *inclusive* effective mass is shown, therefore the $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ requirement cannot be illustrated properly. As devised, nearly all SM background events are removed by these requirements, while the effect on the high mass scenario is not as severe. Figure 5.8 (c) shows the difficulty of compressed scenarios, with very few events in the high m_{eff} region, which requires a lower m_{eff} cut.

5.3.2 Optimisation for compressed spectra

In order to achieve the highest sensitivity of the analysis for our preferred SUSY signals, the precise signal region selections need to be optimised. The SRs used in the previous $\sqrt{s} = 7$ TeV analysis [169], given in table 5.2, cannot be automatically reused: the centre-of-mass energy has increased, the instantaneous luminosity and thus pile-up has increased, and the recorded integrated luminosity is higher. The increased centre-of-mass energy leads to harder objects and higher cross sections for heavy particle production. Our highest signal sensitivity is in the tail of the m_{eff} and other distributions, and with a higher total number of events we have the possibility to set cuts at higher values of the discriminating variables.

Yet this optimisation should also be done in view of the $\sqrt{s} = 7$ TeV results. Fig-

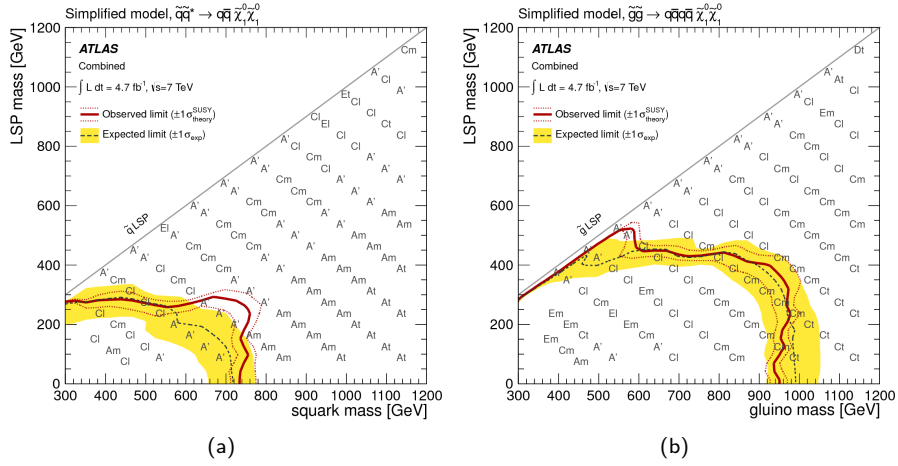


Figure 5.9: The 95% CL exclusion limits with 4.7 fb^{-1} of $\sqrt{s} = 7$ TeV data, for the simplified models with (a) a squark-pair and (b) a gluino-pair decaying directly to quarks and LSPs [169]. The expected exclusion limit is indicated by the black dashed line, with its experimental uncertainties given by the yellow band. The observed limit is indicated by the solid red line, while the dotted lines represent the limit when the signal expectation is deviated by its $\pm 1\sigma$ theoretical uncertainty. For each signal point the SR with the best expected limit is labelled, see table 5.2 for details.

Figure 5.9 shows the $\sqrt{s} = 7$ TeV exclusion limit in two simplified direct decay models. In the left figure, the limit on the squark and LSP mass are given for squark-pair production, while in the right the limit is on gluino and LSP mass in the gluino-pair production simplified model. The maroon solid lines give the observed 95% CL exclusion limit for all SRs combined, with its theoretical uncertainty from PDF and scale uncertainties in the dotted lines. The dashed black line illustrates the expected exclusion limit, with the total experimental uncertainty given by the yellow band. Note that these are combined limits, where for each signal point the results of the SR with highest expected sensitivity is used for both expected and observed exclusion. These best expected SRs per point are indicated by the labels at each point.

An interesting feature is seen. As the production cross section does not depend on the LSP mass (see figure 5.5), one might naively expect to exclude independently of the LSP mass. Yet the limit seems to reach a plateau for a certain LSP mass, meaning that the mass difference between the squark or gluino and the LSP affects the kinematics of the event greatly, and thus the selection efficiency as well.

In view of this observation, the sensitivity of the ‘0-lepton’ analysis has been optimised for the 2012 run, both for scenarios with a high mass for the squark or gluino, and for scenarios with an intermediate $m_{\tilde{q}}$ or $m_{\tilde{g}}$ but a small mass difference

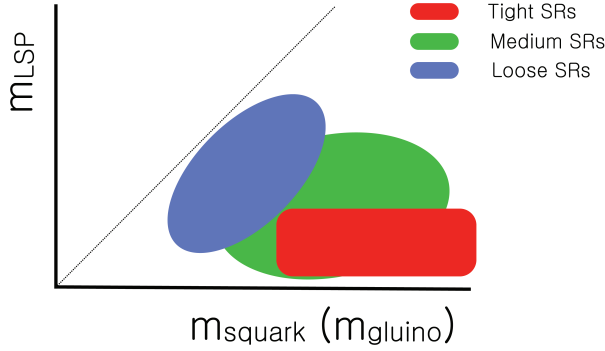


Figure 5.10: Illustration of the areas in our simplified models where the ‘loose’ (light blue), ‘medium’ (green) and ‘tight’ (red) SRs are expected to be most sensitive.

$\Delta m = m_{\tilde{q}/\tilde{g}} - m_{\tilde{\chi}_1^0}$ between the squark/gluino mass and the LSP mass.

Optimisation overview

To be fully consistent with the final analysis, the optimisation of the SRs is done with the exact setup as the final analysis – i.e. with the use of data in control regions to constrain the background prediction, using the statistical procedure described in section 5.5, including all uncertainties described in section 5.4.6. To increase the confidence in the final result, this is preferred over a MC only optimisation, for which the confidence in the validity of the background prediction is less. The optimisation was performed to find the highest sensitivity for up to three selections per jet multiplicity channel. In other words, for each inclusive jet channel, a large set of candidate selections is studied to find those which have the highest expected exclusion limit at 95% Confidence Level (CL) in the mass plane.

Up to three selections are defined: a ‘loose’, ‘medium’ and ‘tight’ selection. The ‘loose’ selections target signal models with a small mass difference between the LSP and gluino or squark, $\Delta m(\tilde{\chi}_1^0, \tilde{g}(\tilde{q})) \lesssim 100$ GeV, while the ‘tight’ selections focus on models with high squark/gluino masses with a large mass difference. The ‘medium’ selections target those in between. This is illustrated in figure 5.10. In the 2011 analysis, a similar set of selections was used, which can be seen from the best expected SRs given as the labels in the exclusion limit plots in figure 5.9. For reference, the precise SR selections of the 2011 $\sqrt{s} = 7$ TeV analysis are given in table 5.2. Note that the ‘loose’ SRA is dubbed A' in figure 5.9. For the squark-pair model, the prevailing SRs are A' , A_m and C_m , which are the ‘loose’ 2-jet and ‘medium’ 2- and 4-jet SRs of the $\sqrt{s} = 7$ TeV analysis. The ‘loose’ SR, which had a looser m_{eff} cut and higher $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ cut, is the most sensitive for very compressed signal points, while the 4-jet SR wins for slightly less compressed scenarios. For the gluino-pair model, the ‘loose’

2-jet SR is again the best performing selection for very compressed scenarios, while 'loose', 'medium' and 'tight' 4-jet SRs are best for less compressed models up to high mass models.

Details of the optimisation procedure

Three direct decay simplified model grids are considered in the optimisation procedure: a model with squark-pair production, one with gluino-pair production and one with associated squark-gluino production. In all of these, the SUSY particles decay into quarks and the LSP directly. In this thesis, the focus is placed on the compressed scenarios, i.e. near the diagonal of the plots. The optimisation for high mass and large Δm was performed in a similar manner on a CMSSM grid.

As mentioned before, we expect that the sensitivity for compressed scenarios will be highest with looser cuts on m_{eff} than for high mass scenarios, as in a hypothetical example with infinite signal statistics, the m_{eff} cut would be placed right on the peak of the signal m_{eff} distribution, which lies roughly at $m_{\text{eff}} \sim 2 \times M_{\text{SUSY}}^{\text{eff}}$, see equation 5.2 and figure 5.3. In these looser SRs, the requirement on the $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ ratio will most probably be tighter, to retain a somewhat stringent requirement on the $\cancel{E}_T$ ³. An important issue is the statistics in both SR and CRs: cutting too tight will lead to low statistics in the CRs, and thus high systematic uncertainties, which will result in a lower sensitivity. A balance needs to be struck between the number of events in both the SRs and corresponding CRs, and the gains from cutting at high values of m_{eff} .

A wide range of variables can be employed to increase the sensitivity of the analysis. Apart from the effective mass and $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ variables, several others have been incorporated in the first phase of the optimisation procedure. Two variables defined for mass reconstruction are the 'transverse mass' m_{T2} [245–247], and the 'contransverse mass' m_{CT} [248, 249]. The former is used for mass reconstruction of pair-produced particles decaying into two invisible particles in the final state, while m_{CT} is a Lorentz invariant variable similar to the transverse mass, but suitable in situations with two particles which have boosts of equal magnitude in opposite directions.

Apart from mass reconstructing variables, also the significance of the missing transverse momentum $\cancel{E}_T^{\text{sig}} \equiv \cancel{E}_T/\sqrt{H_T}$ has been used, with H_T the scalar sum of the transverse momenta of all jets, $H_T = \sum_j p_T^j$. However, none of these variables gives a significantly higher sensitivity for compressed spectra. Therefore, the final optimisation was performed using only the ratio $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ and the effective mass m_{eff} . Furthermore, lowering the p_T threshold of the leading jet to 100 GeV (the lowest value possible with the current trigger at reasonable, but not full, efficiency) does not improve the sensitivity for any of the studied models.

³As a (soon-to-be-seen real-life) example, take a 2-jet SRA 'tight' with cuts $\cancel{E}_T/m_{\text{eff}}^{\text{excl}} > 0.3$ and $m_{\text{eff}} > 1900$ GeV, and a 2-jet SRA 'loose' with $\cancel{E}_T/m_{\text{eff}}^{\text{excl}} > 0.4$ and $m_{\text{eff}} > 1300$ GeV. If no extra jets are available in the event, the former places an effective cut on the missing transverse energy of $\cancel{E}_T > 570$ GeV, while the latter has $\cancel{E}_T > 520$ GeV.

To find the best SRs for each of the 5 jet multiplicities, an ‘exclusion’ fit was performed (see section 5.5) on a blinded dataset of 5.8 fb^{-1} , by varying the $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ cut over the range 0.2–0.5, in steps of 0.025, and the m_{eff} cut between values in the range 800–2000 GeV, in steps of 100 GeV. The blinding of the dataset implies that the measured data in the SRs is not used or even looked at, while the data in CRs are used to estimate the SR backgrounds. For each point on the SUSY model grids the expected excluded signal cross section is obtained through the fit, together with the expected CL_s . This fit is then interpreted in the 2-dimensional $(m_{\tilde{q}}, m_{\tilde{\chi}_1^0})$ or $(m_{\tilde{g}}, m_{\tilde{\chi}_1^0})$ phase space as a 95% exclusion limit, interpolated between the grid points. As the fitter and analysis were still evolving during the optimisation process, the fit was not converging for a few combinations of cuts on certain signal points. This has the effect that the limits are difficult to evaluate on parts of the very compressed regions, however there is sufficient information to come to a conclusion for the final SR selection.

As mentioned before, to limit the total number of SRs of the analysis, a compromise needs to be struck – for each inclusive jet multiplicity channel, only one or two signal selections (‘loose’ or ‘medium’) are defined which perform the best over all three simplified direct decay models. In the following, the results of this optimisation are shown.

Directly decaying squark-pair model

For the $\tilde{q}\tilde{q} \rightarrow q\bar{q}\tilde{\chi}_1^0\tilde{\chi}_1^0$ model, the $\sqrt{s} = 7 \text{ TeV}$ analysis placed the exclusion limit on the LSP mass at $m_{\tilde{\chi}_1^0} \lesssim 300 \text{ GeV}$ for $m_{\tilde{q}} \lesssim 740 \text{ GeV}$. This leaves a large ‘gap’ in sensitivity for higher LSP mass, therefore effort has been put in the optimisation of this region. SRs A, B and C (the 2, 3 and 4 inclusive jet channels respectively) are expected to give the best sensitivity in this model, as a squark-pair decays into two jets, plus possible additional jets from ISR and FSR.

For highly compressed SUSY spectra, additional ISR jets are essential. In these scenarios, due to the small mass difference between the squark and the LSP it decays to, there is less momentum available for the other decay products, and thus it is less likely that one of the final state quarks forms a jet which passes the trigger requirements.

A set of best performing SRs for this model is interpreted in the squark-pair simplified model in figure 5.11. The limits are only shown for the 2, 3, and 4 jet channels, as the higher jet multiplicities do not have a better sensitivity.

The expected exclusion limits do not go higher in LSP mass than the $\sqrt{s} = 7 \text{ TeV}$ results shown in figure 5.9 (a), even though the SR selections from the 2011 analysis have been included in the optimisation procedure. Considering the increase in centre-of-mass energy, together with a larger dataset (5.8 fb^{-1} against 4.7 fb^{-1} in 2011), one would naively hope that the sensitivity would improve. Yet in 2011, it already was not an issue of a too small cross section: the cross section only decreases as a function of squark mass, not LSP mass – therefore, a small increase in statistics will not improve the sensitivity. The problem for these models is actually the signal acceptance and precise simulation of it: for increasing LSP mass, and decreasing

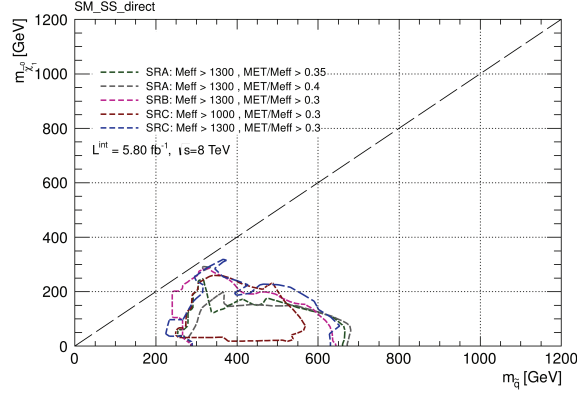


Figure 5.11: The best expected 95% CL exclusion limits for the squark-pair simplified model. The various selections in the 2-, 3- and 4-jet channel are shown in different colours, while the 5- and 6-jet channels are not shown as they add no sensitivity to the compressed region.

Δm , the signal behaves more background-like, as already illustrated by figure 5.2 and seen in the distributions of figures 5.1 and 5.4. In these regions, the systematic uncertainty on background and signal is therefore the limiting factor. As will be discussed in section 5.6, the systematic uncertainty on background for the low jet multiplicity channels is of the order of 20%, mostly driven by a large uncertainty on $W + \text{jets}$. The large ISR uncertainty on signal, up to 40% for models with low mass and low mass splitting, decreases the sensitivity even more. Interestingly, the three and four jet selection outperform the two jet selection for more compressed signal models, where additional ISR and FSR jets will accompany the 2 jets from the squark decays, and there is significantly less background. The best performing selections per jet multiplicity from this optimisation are compared to those from the optimisation of the gluino-pair and squark-gluino production models, after which the best overall are picked.

When the final results are shown in section 5.6, we will see that the high mass region is less affected by the large uncertainties, increasing the excluded squark mass with ~ 50 GeV with respect to the 2011 analysis. In chapter 6 we will discuss how to overcome the limitations of the high systematic uncertainty for the compressed spectra.

Directly decaying gluino-pair model

The exclusion limit from 2011 for the gluino-pair simplified model (figure 5.9 (b)) shows the same plateau-like behaviour as the squark-pair model at $m_{\tilde{\chi}_1^0} \approx 450$ GeV, although one SR containing less events than expected causes a bump in the observed limit along the diagonal.

Figures 5.12 (a) and (b) show the set of best overall SRs interpreted in this model.

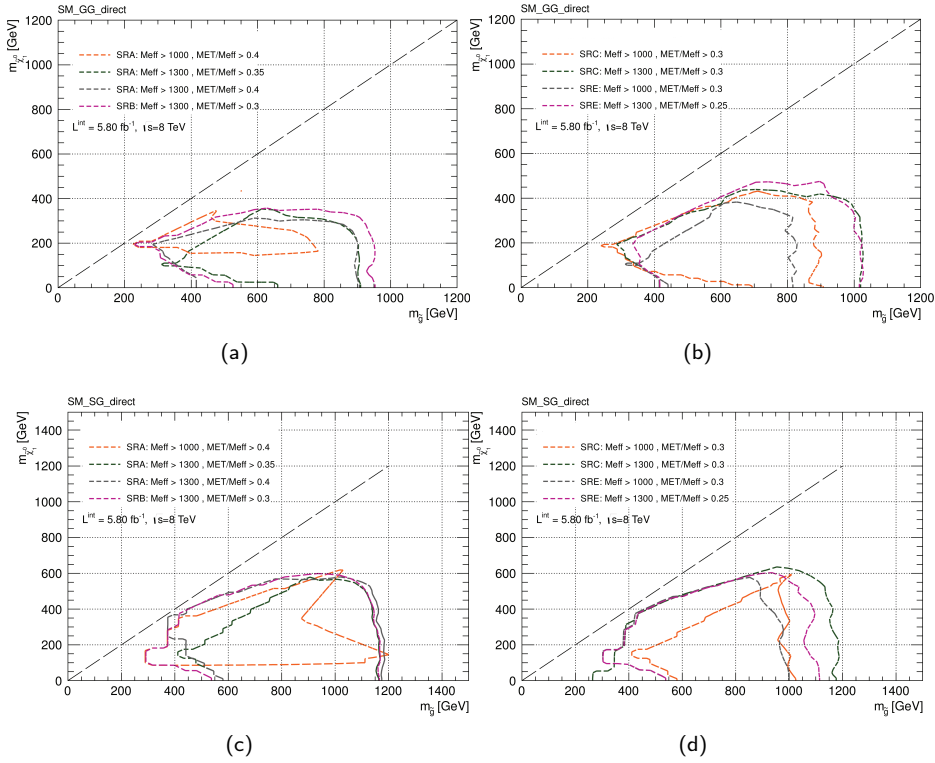


Figure 5.12: The best expected 95% CL exclusion limits for the gluino-pair simplified model (top row) and squark-gluino simplified model (bottom row). The 2- and 3-jet channels are given on the left (a and c), while the 4-, 5- and 6-jet channels are given on the right (b and d). The various selections are shown in different colours. The rectangular feature for SRA with $m_{\text{eff}} > 1000$ GeV and $\cancel{E}_T/m_{\text{eff}}^{\text{excl}} > 0.4$ in figure (c) is caused by non-converging fits for some signal samples.

For a small part in the compressed region, the loosest 2-jet selection with $\cancel{E}_T/m_{\text{eff}}^{\text{excl}} > 0.4$ and $m_{\text{eff}} > 1000$ comes closest to the diagonal, as is seen by the orange line in figure 5.12 (a). For higher gluino masses, the 4- and 6-jet regions perform significantly better, limiting the LSP mass to be higher than ~ 500 GeV.

Directly decaying squark-gluino model

Using the same procedure, the focus is placed again on the phase-space region with intermediate gluino and squark mass, and LSP mass close to these masses. As $\tilde{q}\tilde{g}$ production has the highest production cross section of these three simplified models, we can naively expect the exclusion limit to reach the farthest in terms of gluino mass.

From the kinematics of the process, we expect to have at least 3 jets in the signal events, as the $\tilde{q}$ and $\tilde{g}$ decay to 1 and 2 quarks, respectively, each forming a jet. Additional radiation from either ISR or FSR can lead to signatures with 4-5 jets, while jets can be misreconstructed or fall out of the acceptance, decreasing the number of jets.

The best performing SR selections are shown for the 2- and 3-jet channels in figure 5.12 (c) and for 4- and 6-jet channel in figure 5.12 (d). None of the SRs comes close to the diagonal when $m_{\tilde{g}} > 500$ GeV, while several follow the same slope. As many of the combinations of m_{eff} and $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ close by those combinations shown in the figure lead to equivalent behaviour in the compressed region, and the exclusion limit in the high-mass region will be driven by the ‘tight’ regions, the final selections of the ‘loose’ and ‘medium’ regions are driven by the optimisation in the squark-pair and gluino-pair production models.

5.3.3 High mass optimisation

To optimise the analysis for SUSY scenarios with gluinos and/or squarks with high mass and low-mass LSPs, ‘tight’ signal regions are defined using a similar procedure as for the compressed scenarios discussed above. Exclusion limits, derived with the full statistical analysis using data in CRs, are compared for a range of values of m_{eff} and $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ in the CMSSM. $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ is varied between 0.15 and 0.35, while m_{eff} is varied between 1200 GeV and 2200 GeV, with steps of 100 GeV. The jet p_T requirement for all subleading jets is studied as well. The second hardest jet p_T requirement is varied between 60 GeV and 100 GeV, where 60 GeV is indeed the best performing value. Likewise, although softer fifth and sixth jets (40 GeV) in SRD and SRE increase signal acceptance, the background rejection decreases significantly, leading to worse exclusion sensitivity.

As in the optimisation for compressed SUSY, low jet multiplicity analysis channels have tighter requirements on both $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ and m_{eff} .

5.3.4 Final SR selection

The final SR selection is given in table 5.3. The SRs which have been optimised for high sensitivity for compressed SUSY are denoted by ‘loose’ and ‘medium’. The optimisation of the ‘tight’ SRs, targeting high mass scenarios indeed lead to high m_{eff} cuts, while the $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ requirements are lower than for the ‘loose’ and ‘medium’ SRs as expected. Note that not all jet multiplicities have all three types of SR: to reduce the number of total SRs, regions which have overlapping sensitivity have been removed. In total 12 SRs are defined, spread over the 5 jet multiplicities. The final results are shown and discussed in section 5.6. To show that the ‘loose’ and ‘medium’ selections are indeed more sensitive in the compressed region, the final expected 95% CL exclusion limits for each of the SRs are shown in that section as well for all three simplified models.

The relative number of simulated signal events passing the above requirements, known as the signal efficiency, is shown for directly decaying squark-pair models in

Cut Description	A (l/m/t)	B (m/t)	C (l/m/t)	D (t)	E (l/m/t)
Data quality	Good Run List				
Trigger	EF_j80_a4tchad_xe100_tclcw_veryloose				
Cleaning cuts	See text				
Lepton veto	No e (μ) with $p_T > 20$ (10) GeV after overlap removal				
$\cancel{E}_T$ [GeV] >	160				
Jet 1 p_T [GeV] >	130				
Jet multiplicity	≥ 2	≥ 3	≥ 4	≥ 5	≥ 6
Jet 2 p_T [GeV] >	60	60	60	60	60
Jet 3 p_T [GeV] >	–	60	60	60	60
Jet 4 p_T [GeV] >	–	–	60	60	60
Jet 5 p_T [GeV] >	–	–	–	60	60
Jet 6 p_T [GeV] >	–	–	–	–	60
$\Delta\phi(j_i, \cancel{E}_T)$ > ($i \leq 3$)	0.4	0.4	0.4	0.4	0.4
$\Delta\phi(j_i, \cancel{E}_T)$ > ($i > 3$)	–	–	0.2	0.2	0.2
$\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ >	0.4/0.4/0.3	0.3/0.25	0.3/0.3/0.25	0.15	0.3/0.25/0.15
$m_{\text{eff}} >$ [GeV]	1000/1300/1900	1300/1900	1000/1300/1900	1700	1000/1300/1400

Table 5.3: The event selection for the $5.8 \text{ fb}^{-1} \sqrt{s} = 8 \text{ TeV}$ 0-lepton analysis. The separation in five analysis channels is performed from the jet multiplicity cut onwards, while the separation between ‘loose’ (l), ‘medium’ (m) and ‘tight’ (t) is done in the $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ and m_{eff} cuts. For each inclusive jet multiplicity channel, additional jets are required to have $p_T > 40 \text{ GeV}$.

figure 5.13 and for gluino-pair models in figure 5.14, for events with at least 2-jets, and each of the loosest SRs of every jet multiplicity channel. Whereas the production cross sections are independent of the LSP mass (see figure 5.5), the signal efficiency decreases with decreasing $m_{\tilde{q}} - m_{\tilde{\chi}_1^0}$. This is exactly what we expected, and will have a large effect on any interpretation of the results.

5.4 Background estimation

One of the most critical parts of any analysis is background estimation. To evaluate the number of possible signal events in our SRs, one needs to know how many background events have found their way inside the selection.

To estimate the SM background in each of the signal regions, the shape of the SM

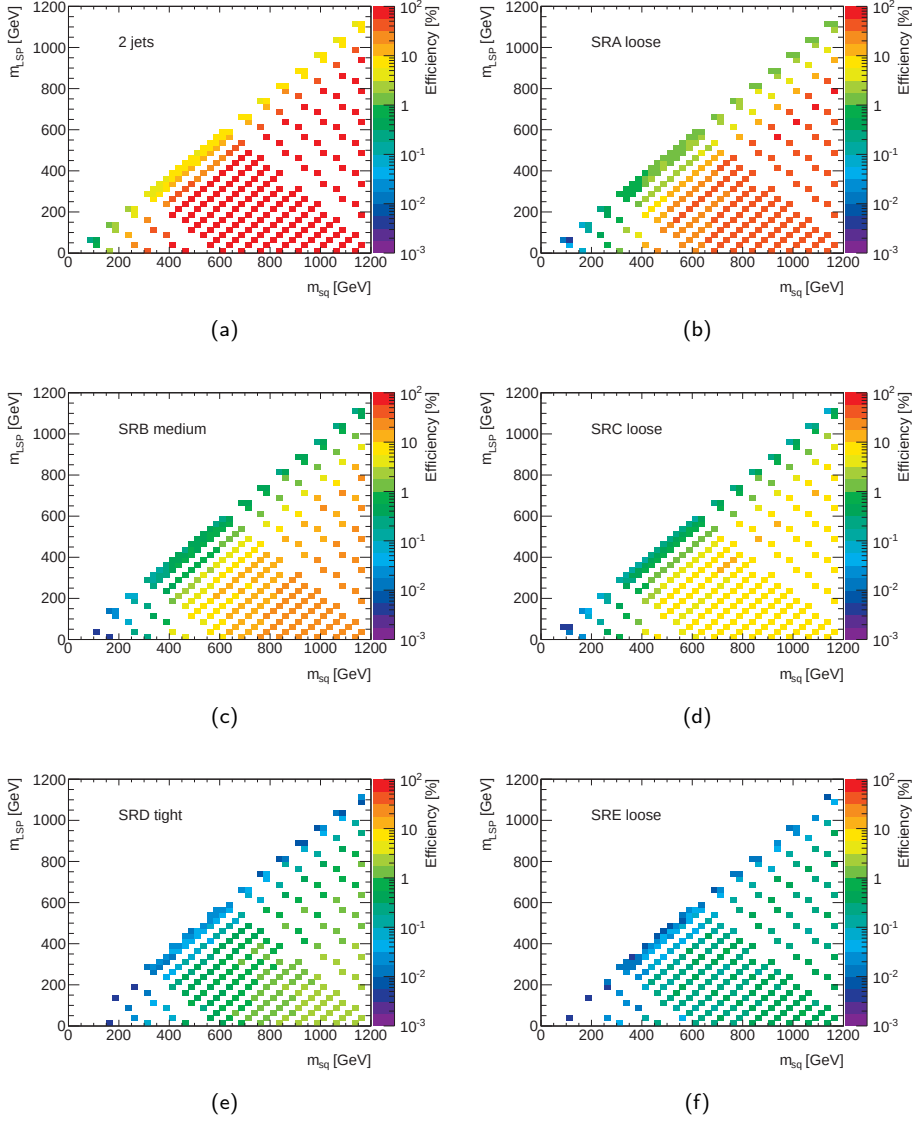


Figure 5.13: Signal efficiency for the squark-pair simplified model after the $\cancel{E}_T$ and 2 jet selection plus $\Delta\phi(jet, \cancel{E}_T)$ requirements (a), and after the full SR selections for SRA 'loose' (b), SRB 'medium' (c), SRC 'loose' (d), SRD 'tight' (e) and SRE 'loose' (f).

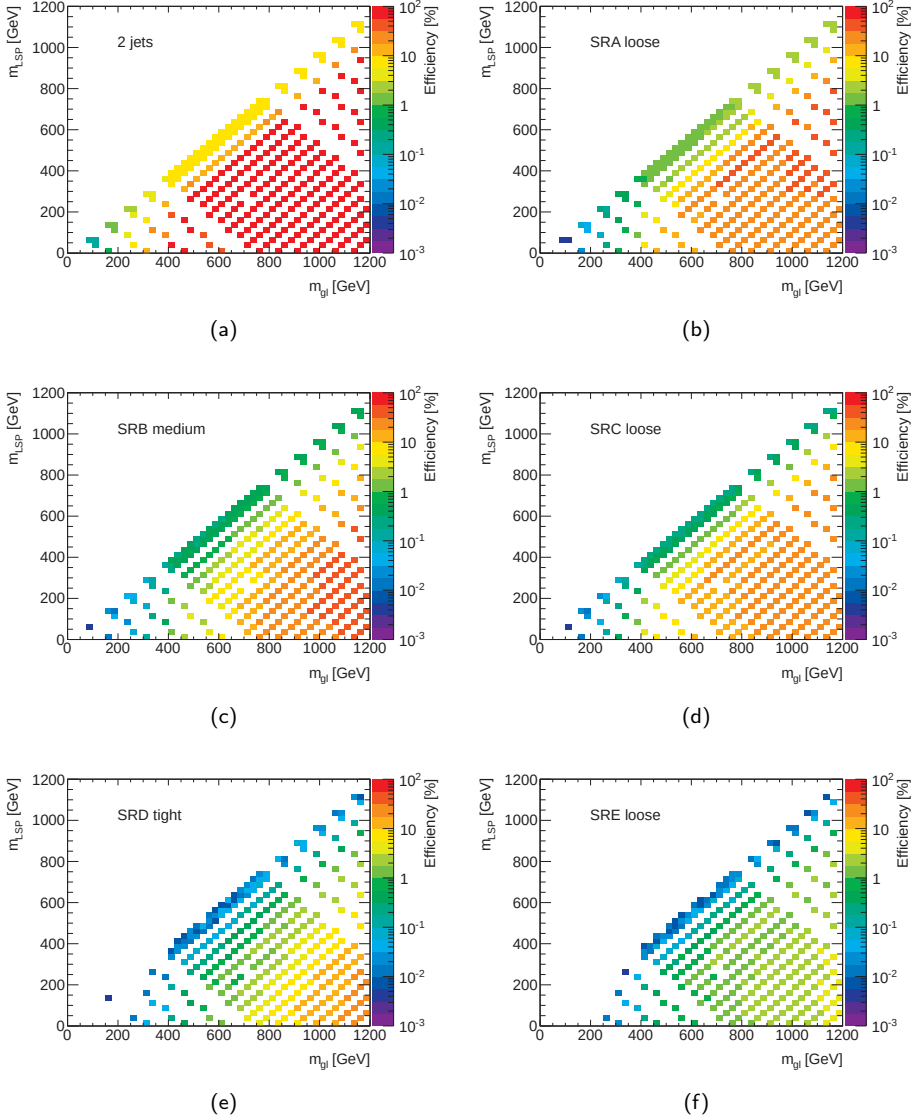


Figure 5.14: Signal efficiency for the gluino-pair simplified model after the $\cancel{E}_T$ and 2 jet selection plus $\Delta\phi(\text{jet}, \cancel{E}_T)$ requirements (a), and after the full SR selections for SRA ‘loose’ (b), SRB ‘medium’ (c), SRC ‘loose’ (d), SRD ‘tight’ (e) and SRE ‘loose’ (f).

background is taken from MC simulation, while the normalisation of each background is obtained from the control regions through the statistical procedure explained in section 5.5. Four background processes have their own CR, whose selections are designed to be enriched in that particular process: QCD multi-jets, $Z(\rightarrow \nu\nu)+\text{jets}$, $W^\pm(\rightarrow l^\pm\nu)+\text{jets}$ and semileptonic $t\bar{t}$. These CRs are defined as kinematically close to the SRs as possible, ensuring that events falling in the CRs will be very similar to the SM background events ending up in our SRs. Thus an inclusive 2-jet SR will have inclusive 2-jet CRs with similar m_{eff} cuts. The use of control regions is preferred over the direct use of MC simulated data, to be independent on the MC normalisation of the events. By using this semi-data driven method uncertainties on MC modelling are reduced to uncertainties on the shapes of the distributions.

5.4.1 Control regions

Events coming from $W^\pm(\rightarrow l^\pm\nu)+\text{jets}$ or semileptonic $t\bar{t}$ decays ($t\bar{t} \rightarrow W^\pm q W^\mp q \rightarrow l^\pm \nu qq' q'$) have a similar signature, and can contaminate the SR if the electron or muon is outside of the detector acceptance or misidentified. To define the CRs, events coming from these processes are preselected by requiring precisely one electron or muon and $E_T > 160$ GeV. Furthermore, in order to reconstruct the leptonically decaying W boson, the transverse mass constructed from the lepton and the neutrino, represented by $\cancel{E}_T$, should be consistent with such a W decay: $30 < m_T(\cancel{E}_T, l) < 100$ GeV. $t\bar{t}$ decays are distinguished from W decays by the presence of a jet from a b -quark. The top control region (CR_T) therefore requires a b -jet. The control region enriched in W production (CR_W) on the other hand vetoes events with a b -jet. For both $t\bar{t}$ and $W+\text{jets}$, typically $\sim 90\%$ of the events which end up in the SR have a lepton which has been observed as a jet; mostly hadronically decaying τ -leptons and misidentified electrons. Therefore it makes sense to treat the lepton as a jet in the final variables – after the m_T requirement, all jet-dependent variables, such as the number of jets, m_{eff} , $\Delta\phi(j_i, \cancel{E}_T)$, are recalculated. Finally, the jet multiplicity cuts are applied as in the SRs, splitting the events in 5 inclusive jet multiplicity channels.

$Z(\rightarrow \nu\nu)+\text{jets}$ is an irreducible background, implying that it cannot be distinguished from the targeted SUSY signatures with our observables. Thus the approach used for CR_W and CR_T , where the selection targets the precise underlying process, is not viable for $Z(\rightarrow \nu\nu)+\text{jets}$. To estimate the number of these events, a different but similar process to $Z(\rightarrow \nu\nu)+\text{jets}$ needs to be selected. Two such processes are the decay of the Z boson into leptons ($Z \rightarrow l^\pm l^\mp + \text{jets}$) and the production of a photon with jets. The former is kinematically equivalent to the decay to two neutrinos. Selecting two same flavour light leptons of opposite charge, e^+e^- or $\mu^+\mu^-$, and requiring that the invariant mass of the leptons is comparable to the mass of the Z boson, $66 \text{ GeV} < m_{ll} < 116 \text{ GeV}$, results in a sample of predominantly $Z \rightarrow l^\pm l^\mp + \text{jets}$ events. Events need also to pass the jet and $\cancel{E}_T$ requirements of the SR selection, yet to mimic the behaviour of the decay into neutrinos, the p_T of the reconstructed Z boson is added to the $\cancel{E}_T$ vector. A problem with the selection of leptonically decaying Z bosons is that the branching ratio for $Z \rightarrow l^\pm l^\mp$, with $l = e, \mu$, is only

7%, compared to 20% of the decay into neutrinos [19]. The lower number of available events results in low statistics in the selection: in the order of 10 events in the 4-jet channel.

An alternative is given by a γ +jets selection. As both the Z boson and the photon are mixtures of neutral gauge bosons (see chapter 1), interactions involving photons are similar to those with Z bosons. To pass the SR $\cancel{E}_T$ requirement, $Z(\rightarrow \nu\nu)$ +jets events need a source of high missing transverse momentum – events passing the SR requirements typically have the Z boson recoiling against the jets with high p_T . Yet for $p_T \gg m_Z$, the kinematics of the event no longer depends on the mass of the Z boson, and γ +jets events are expected [250, 251] to be kinematically similar to $Z(\rightarrow \nu\nu)$ +jets events. By selecting a sample of photon plus jets events in data and scaling these to the $Z(\rightarrow \nu\nu)$ +jets cross section, a sample with pseudo $Z(\rightarrow \nu\nu)$ +jets events is obtained. To mimic the $\cancel{E}_T$ from the two neutrinos, the photon p_T should be added to the measured $\cancel{E}_T$ in the event. Again, the event should also pass the jet p_T requirements, as well as have (modified) $\cancel{E}_T > 160$ GeV. The photon identification is discussed in section 4.4. Care must be taken that the selected photon is a prompt photon, as photons can also be radiated off charged final state particles, leading to soft or non-isolated photons in the event. The cross section ratio

$$R_{Z/\gamma}(p_T) = \frac{\partial\sigma_{(Z(\rightarrow \nu\nu)+jets)}/\partial p_T}{\partial\sigma_{(\gamma+jets)}/\partial p_T} \quad (5.6)$$

is obtained from MC due to limited statistics in data. The ratio is of the order of 0.3, while the p_T dependence is fairly flat. The cross section ratio is corrected for the $Z \rightarrow \nu\nu$ branching ratio before scaling the selected photon events. Since the γ +jets cross section is significantly larger than that of $Z \rightarrow ll$, and the photon is easy to reconstruct, the dilepton selection defined previously is only used as a validation of the γ +jets method. It constitutes the validation region VR_Z . The photon selection is promoted to being the control region (CR_γ) of choice for the estimation of Z +jets events.

In the previous discussion, contributions to the CRs coming from events with fake leptons have been ignored. It has been shown in the $\sqrt{s} = 7$ TeV analysis [252] that these fake lepton contributions fall rapidly with increasing $\cancel{E}_T$, and amount to no more than 4% in CR_W and even less in CR_T and VR_Z . Therefore, these contributions are neglected.

However, the contribution of QCD multi-jet events with fake photons to CR_γ has been taken into account via the ‘matrix method’ for photons. Fake photons originate from jets containing neutral pions decaying into photon pairs. Fake photons therefore are less isolated than prompt photons, since other particles in the jet will deposit energy close by the detected photon. The matrix method uses this information. Four regions are defined: a signal region A of ‘tight’ photons with calorimetric isolation; an isolated control region C with isolated non-‘tight’ photons; and two control regions B and D with non-isolated ‘tight’ and non-‘tight’ photons. Assuming the signal region has low contamination and the background shape is uncorrelated between the regions,

Control region	Selected process in CR	cuts
CR _W	$W^\pm(\rightarrow l^\pm\nu)+\text{jets}$	$n_e == 1$ or $n_\mu == 1$, b -jet veto $30 < m_T(\cancel{E}_T, l) < 100$ GeV No $\Delta\phi(j_i, \cancel{E}_T)$ and $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ cuts
CR _T	$t\bar{t} \rightarrow l^\pm\nu bqq'b$	$n_e == 1$ or $n_\mu == 1$, b -jet tag $30 \text{ GeV} < m_T(\cancel{E}_T, l) < 100 \text{ GeV}$ No $\Delta\phi(j_i, \cancel{E}_T)$ and $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ cuts
CR _{γ}	$\gamma + \text{jets}$	Isolated photon
CR _{QCD}	QCD multi-jets	reversed $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ cut and $\Delta\phi(j_i, \cancel{E}_T) < 0.2$: for SRA & SRB, $i = 1, 2, (3)$; for SRC-SRE, $i = \text{any jet with } p_T > 40 \text{ GeV}$

Table 5.4: Requirements for each of the four CRs. The second column gives the process which is selected by these cuts. Otherwise, the SR requirements are applied.

the number of fake photons in the signal region can be estimated from the number of observed events in the control regions: $N_A = N_C^{\text{obs}} \cdot N_B^{\text{obs}} / N_D^{\text{obs}}$ [253]. Using this matrix method, the fake contribution is estimated to be less than 5% of the total SM background in CR _{γ} .

Note that in the final selections of CR_W, CR_T and VR_Z, the cuts on $\min(\Delta\phi(j_i, \cancel{E}_T))$ and $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ are not applied. Applying them would lead to very small event counts (less than 5) in the CRs, and thus large uncertainties. Removing these cuts does not affect the ratio between CR and SR. The main CR selections are listed in table 5.4, together with the process which is selected.

The estimated number of QCD multi-jet events in the signal regions is derived via a data-driven jet-smearing technique described in ref. [254]. First, a jet response function is measured from di-jet events in collision data. The jet response function describes the jet energy measurement by the detector, and the corresponding $\cancel{E}_T$ from jet mismeasurements, jet resolution and neutrinos from heavy-flavour decay in jets. With this knowledge on the detector response for the full $\cancel{E}_T$ range, low $\cancel{E}_T$ events are selected in data and convoluted with this response function, resulting in high $\cancel{E}_T$ ‘pseudo events’. These are then used to describe the shape of the distributions in the control and signal regions. To estimate the normalisation of QCD multi-jets in the SR, a specific control region CR_{QCD} is defined, which selects a sample of events enriched in QCD (pseudo) events. From the discussion in section 5.3.1 we know that QCD multi-jet events with a misidentified jet faking real $\cancel{E}_T$ have their $\cancel{E}_T$ pointing in the direction of the jet. The obvious choice for CR_{QCD} is therefore to reverse (and tighten) the requirement on the angle between the jet and the missing transverse momentum: $\min(\Delta\phi(j_i, \cancel{E}_T)) < 0.2$ for all three leading jets with $p_T > 60$ GeV in

SR A and SR B, or all jets with $p_T > 40$ GeV for SR C, D and E. Furthermore, W +jets and semileptonically decaying $t\bar{t}$ events with a boosted W boson decaying to a hadronically decaying τ can have their $\cancel{E}_T$ aligned with a jet as well. Therefore the requirement on $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ is reversed to remove some of these events and improve the purity of the QCD multi-jet selection.

5.4.2 Validation regions

Finally, a set of validation regions (VRs) is defined, of which one (VR_Z) has already been introduced. These validation regions are used to validate the results of the background estimation. Using the same procedure as for the SRs, the CR observations are extrapolated to the VRs using transfer factors. Each of the VRs tests the prediction done by the control regions of a specific background process. Apart from VR_Z , the validation regions are defined for validating W +jets, $t\bar{t}$ and QCD multi-jet backgrounds, and they are defined as:

- VR_{T2L} : a selection focussing on double leptonic $t\bar{t}$ production, by requiring 2 opposite sign leptons with high p_T and $m_{ll} > 116$ GeV;
- VR_{TM} and VR_{WM} : selections as alternative for CR_W and CR_T , where the lepton is treated as a missing particle instead of a jet;
- VR_{Wf} and VR_{Tf} : analogous to CR_W and CR_T but with application of the final $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ and m_{eff} requirements in order to test the assumption that we are allowed to relax these requirements in the CRs.
- VR_{WT+} and VR_{WT-} : analogous to $\text{CR}_W + \text{CR}_T$, but selecting events with positive and negatively charged leptons, respectively.
- VR_{Q1} and VR_{Q2} : selections to check QCD multi-jet estimations. Analogous to SR selections except for $\min(\Delta\phi(j_i, \cancel{E}_T)) < 0.2$ for the leading 3 jets with $p_T > 40$ GeV in VR_{Q1} , to enhance the QCD multi-jet production component and test its extrapolation along $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$. In VR_{Q2} the SR $\min(\Delta\phi(j_i, \cancel{E}_T))$ cut is reversed to check the extrapolation along this variable.

5.4.3 Agreement of data and MC in CRs

To gain confidence in the understanding of the control regions before looking at their final m_{eff} distributions, we study the comparison between data and MC in several interesting distributions. Figure 5.15 shows several distributions in data and MC simulations for the various control regions. The preselection for these distributions is $\cancel{E}_T > 160$ GeV and two jets with $p_T > 130$ GeV for the leading and $p_T > 60$ GeV for the subleading jet. Data points are illustrated by black dots, while the filled histograms denote the raw MC samples. The lower pad in each figure illustrates the ratio of data to MC, together with the systematic uncertainties on the MC simulation. These systematic uncertainties will be discussed in the following section.

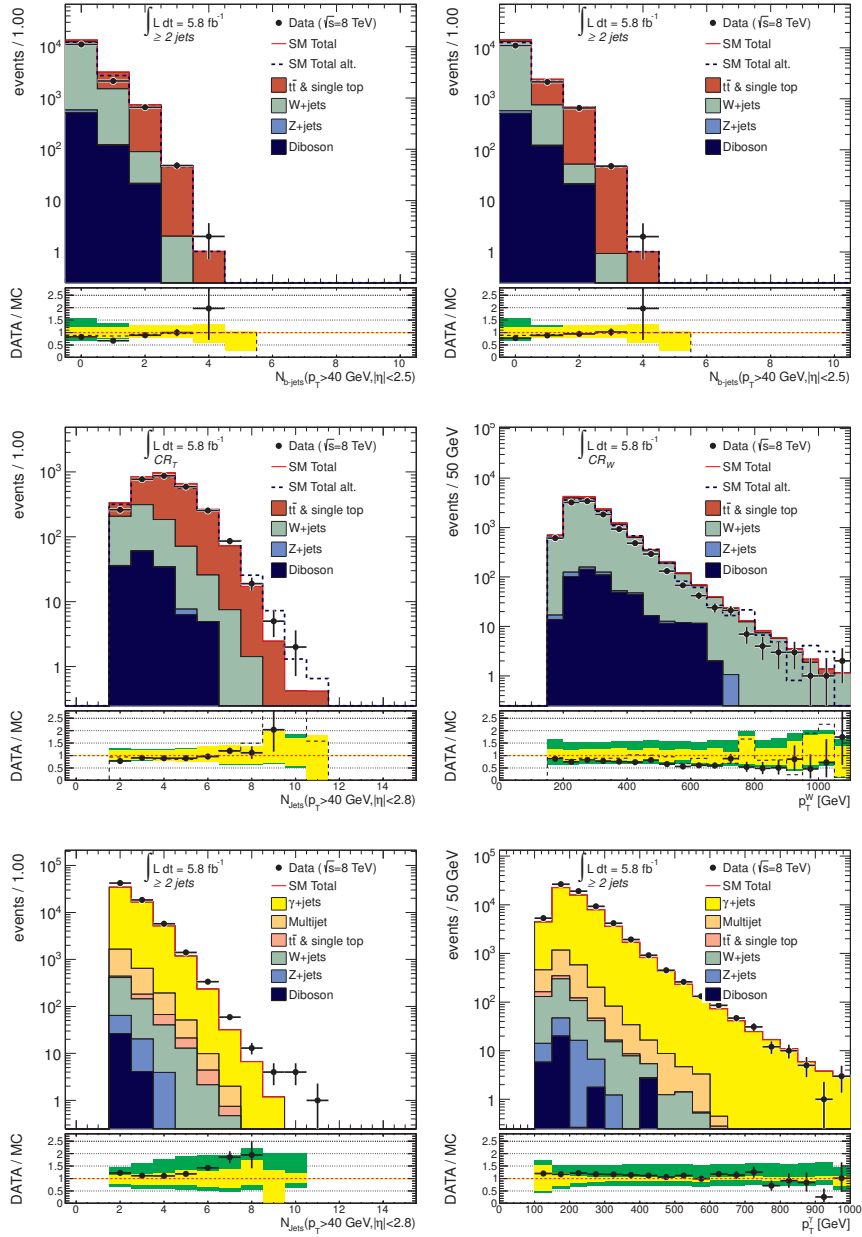


Figure 5.15: Observed and MC distributions of the number of b -jets in $CR_W + CR_T$ (top left), and the same distribution with the SHERPA events rescaled (top right), see section 5.4.4; jet multiplicity in CR_T (middle left); p_T^W of the reconstructed W boson in CR_W (middle right); jet multiplicity (bottom left) and transverse momentum of the photon p_T^γ (bottom right) in CR_γ . In (a), (b) and (d), the ALPGEN W +jets sample is shown as alternative, while in (c) the SHERPA $t\bar{t}$ is the alternative.

In the top left of figure 5.15 the b -jet multiplicity is shown in the combination of $\text{CR}_W + \text{CR}_T$. The recorded data is not well reproduced by the SHERPA sample: in the 0 b -jet bin, data lies $\sim 10\%$ below the MC prediction, while in the 1 b -jet bin the difference is even $\sim 45\%$. This is due to an issue with the implementation of SHERPA, discussed in length shortly, in section 5.4.4. After the rescaling of SHERPA samples in both CR_W and CR_T as described in that section, the issue is resolved. This can be seen from the top right of figure 5.15. This rescaling is therefore performed for all further plots in this chapter. There is however still a $\sim 20\%$ downward fluctuation in the 0 b -jet bin. This automatically means that the MC estimate in CR_W is $\sim 20\%$ too high. No significant discrepancy is seen in the other bins. Alternative MC samples for both W +jets and $t\bar{t}$ events, described in the following, show similar behaviour.

The number of jets in CR_T events shown in figure 5.15 (middle row, left) is reasonably well described by MC@NLO up to 8 jets. The slope of the distribution is slightly too soft in the simulation. The transverse momentum of the reconstructed W boson in CR_W is given in the middle right figure of 5.15. The predicted shape of this p_T^W distribution is correct, yet the 20% offset is clearly visible. CR_W is also seen to be dominated by W +jets events, as devised.

In the bottom row of figure 5.15 the jet multiplicity and photon p_T^γ for events in CR_γ are shown. The latter is again well described, while the number of events in CR_γ with more than 5 jets is underestimated in the used ALPGEN sample, where the slope of the distribution falls off too quickly in the simulation compared to data. The high jet multiplicity CR_γ regions will therefore have an excess of data over the MC prediction.

Alternative MC samples

In order to evaluate how well our MC generators perform, alternative samples have been shown in the distributions for CR_W , CR_T and VR_Z . As different generators have different underlying principles, the modelling uncertainty is evaluated by comparing the ‘nominal’ MC sample to the alternative one.

In the CR_T jet multiplicity distribution (middle left in figure 5.15), the alternative MC sample generated by SHERPA predicts too many jets, and is in other words overcompensating the too soft slope seen in MC@NLO. This is the main reason of preferring MC@NLO over SHERPA for $t\bar{t}$. The difference between the generators is taken as a theoretical uncertainty on $t\bar{t}$ to account for the non-negligible differences between the generators. In the m_{eff} distributions shown and described in the next paragraph the discrepancy will be evident as well.

In CR_W ALPGEN is used as an alternative sample for W +jets events, which shows similar shapes to SHERPA. However, the number of simulated events in the available ALPGEN samples is not sufficient for the ‘tight’ SRs. For Z/γ^* +jets events SHERPA is used as alternative.

5.4.4 Sherpa heavy-flavour issue

For the SHERPA W +jets and $t\bar{t}$ samples, a problem was found in the flavour tagging: in some events (up to 25% for $t\bar{t}$) the vertex of heavy-flavour jets is set to be the origin of the coordinate system due to a bug in the implementation of the algorithm, causing them to be reconstructed as light-flavour jets. The consequence is that the b -jet multiplicity is affected for the W +jets sample, as well as for the alternative $t\bar{t}$ sample which is also generated with SHERPA. This effect is observed in the top left figure in 5.15. Since the statistics of the ALPGEN W +jets sample is not high enough to have a sufficient number of events in the tight regions, the implemented solution is to reweigh the SHERPA sample by comparing the b -jet multiplicity distribution to ALPGEN. To rescale the samples in CR_W , the relative number of events without a tagged b -jet

$$r = \frac{N_{nob-tag}}{N_{total}} \quad (5.7)$$

is calculated for both SHERPA and ALPGEN. The SHERPA events are corrected such that their r agrees with that of ALPGEN: for CR_W , the SHERPA events are thus rescaled by a factor $r_{ALPGEN}/r_{SHERPA} = 0.95/0.89$. In CR_T , the relative number of b -tagged events is compared, leading to a rescaling of SHERPA W +jets by $(1-0.95)/(1-0.85)$. These rescaling factors on the SHERPA W +jets are applied in all distributions shown in the rest of this chapter.

5.4.5 Final CR distributions

The final $\cancel{E}_T/m_{\text{eff}}^{excl}$ and m_{eff} distributions for CR_T , CR_W , CR_γ and CR_{QCD} are shown in figures 5.16, 5.17, 5.18 and 5.19, respectively, for the 2- and 4-jet channels. These distributions have all CR requirements applied to them except the final $\cancel{E}_T/m_{\text{eff}}^{excl}$ and m_{eff} cuts. Note that CR_W and CR_T have no $\Delta\phi$ and $\cancel{E}_T/m_{\text{eff}}^{excl}$ requirements applied. The MC simulation seems to describe the data quite well in all the CRs, although the overshoot of data in the 4-jet channel of CR_γ above the ALPGEN γ +jets sample is already apparent. Also, there are more SHERPA W +jets events surviving in the tail of the CR_W m_{eff} distribution after the ‘tight’ m_{eff} cuts. From these measurements, the final background estimation in the SRs is performed using the transfer factors as given by equation 5.4.

5.4.6 Sources of systematic uncertainties on SM background

The sources of uncertainty on the SM background estimation are varied, but fall mostly in one of three categories:

- Experimental uncertainties, related to the understanding of the detector. The most prominent sources of these uncertainties are: jet energy scale (JES) and the jet energy resolution (JRES), which both affect the jet p_T distribution and the $\cancel{E}_T$ calculated from these jets; scale and resolution of the soft terms of the $\cancel{E}_T$, $(\cancel{E}_T)^{softJets}$ and $(\cancel{E}_T)^{CellOut}$; the b -tagging efficiency in CR_W and CR_T ;

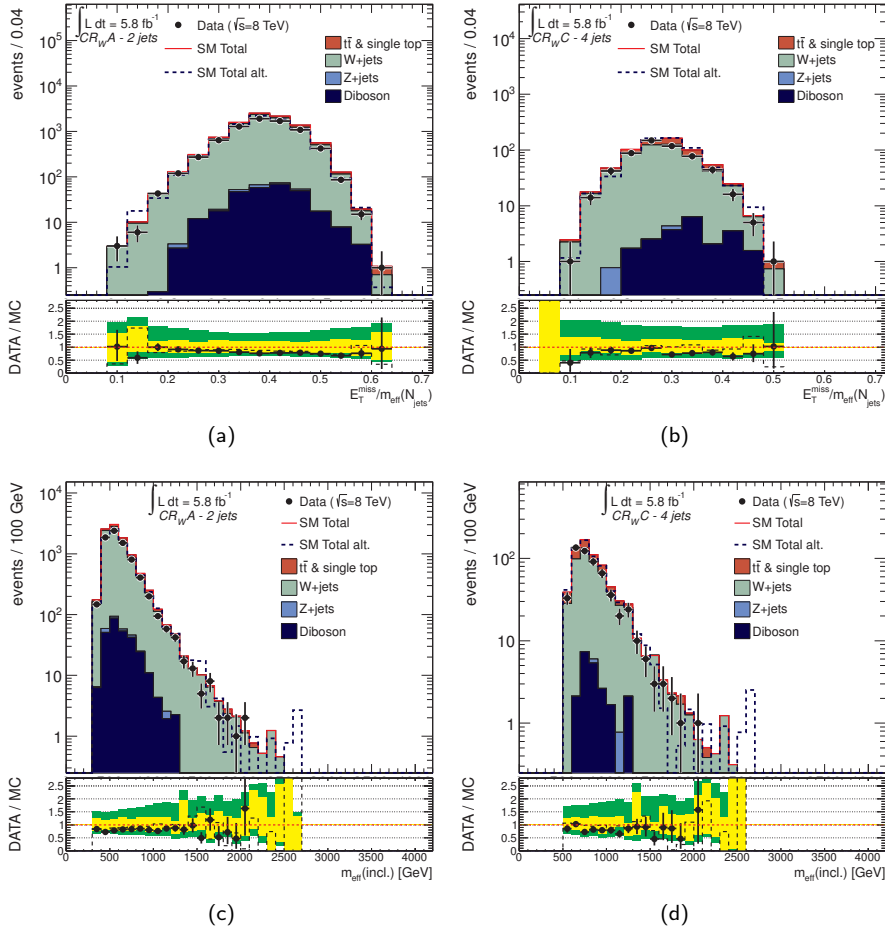


Figure 5.16: Observed and simulated $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ and m_{eff} distributions in the 2- and 4-jet channel for CR_W .

and the pile-up, which affects the jet multiplicity and the $\cancel{E}_T$. The underlying principles of these uncertainties are all described in chapter 4. Also, the measurement of the total integrated luminosity comes with an uncertainty.

- Theoretical uncertainties of the modelling of the MC. The choice of PDF, factorisation and renormalisation scale in the MC generator all have an effect on the final simulated events. Furthermore, the production of ISR jets in signal samples has an uncertainty. Finally, a bug in the heavy-flavour description as described in 5.4.4 leads to an additional uncertainty of the MC modelling with SHERPA.
- Uncertainties on the extrapolation from CR to SR. Statistical uncertainties orig-

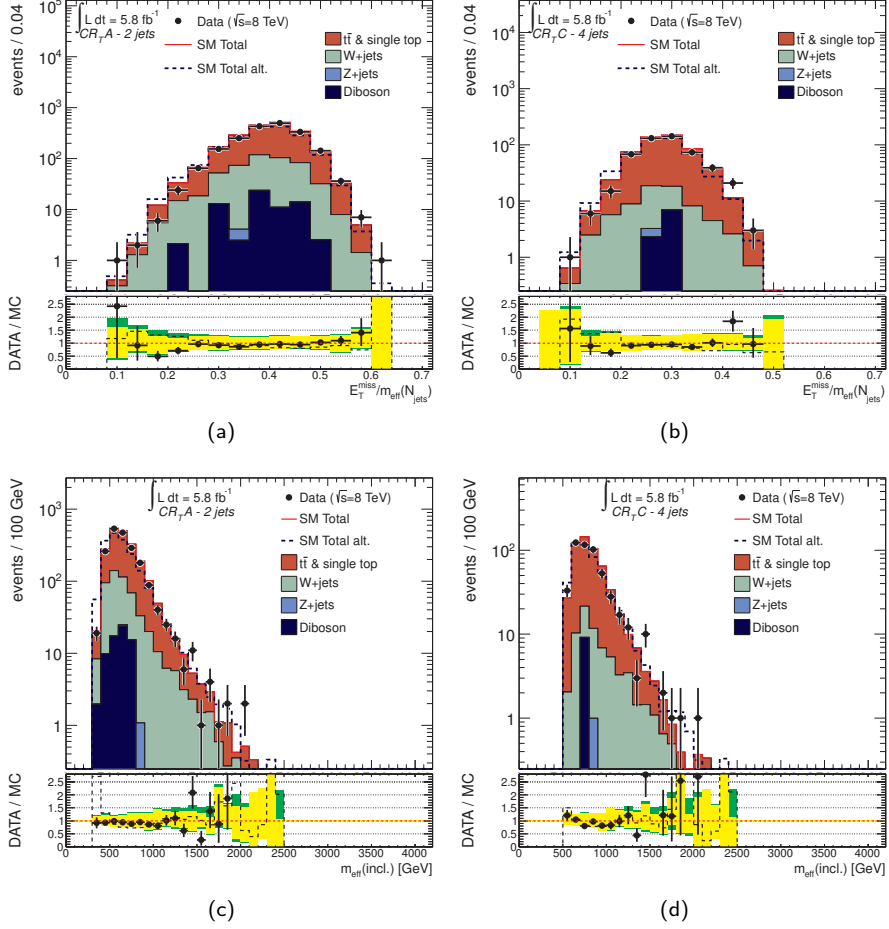


Figure 5.17: Observed and simulated $E_T/m_{\text{eff}}^{\text{excl}}$ and m_{eff} distributions in the 2- and 4-jet channel for CR_T .

inating from the limited size of both data and MC datasets in the CRs propagate to an uncertainty on the estimated number of background events in the SR.

Since the backgrounds in the SRs are predicted from the CR counts using transfer factors, each of these uncertainties propagates to an uncertainty on the transfer factor. The uncertainties coming from detector effects have been studied by dedicated analyses, resulting in p_T and η dependent functions describing these uncertainties for jets, E_T and leptons. For both background and signal MC, uncertainties coming from JES and the soft terms of the E_T are evaluated by comparing the number of SR events after scaling the jet p_T s and E_T up and down by one standard deviation, given by their respective uncertainty functions. To evaluate the uncertainty on the jet resolution, the

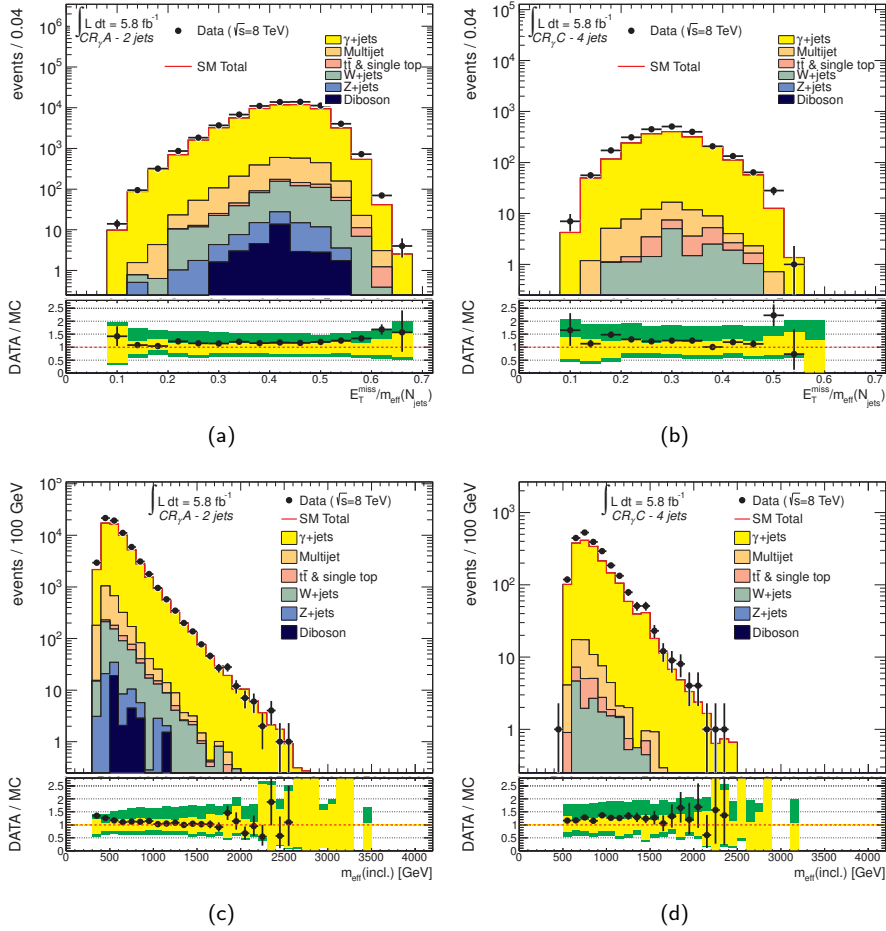


Figure 5.18: Observed and simulated $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ and m_{eff} distributions in the 2- and 4-jet channel for CR_γ .

jet momenta and $\cancel{E}_T$ are smeared by a factor obtained from a Gaussian with a width given by the jet resolution. When determining the impact of these uncertainties on the analysis, the modified objects are propagated in the analysis, affecting $\cancel{E}_T$ and the number and p_T of the jets. In CR_W and CR_T a b -tagging algorithm is used to find heavy-favour jets, and separate W +jets from $t\bar{t}$ events. The uncertainty on the b -tagging efficiency is given by an error on the scale factor from MC to data, see section 4.2.2. These uncertainties are taken into account in both CR_W and CR_T . Finally, uncertainties due to the simulation of pile-up are evaluated by varying the mean number of interactions per event $\langle\mu\rangle$ in MC by a factor of 0.9 and 1.1, which covers the differences seen in figure 5.20.

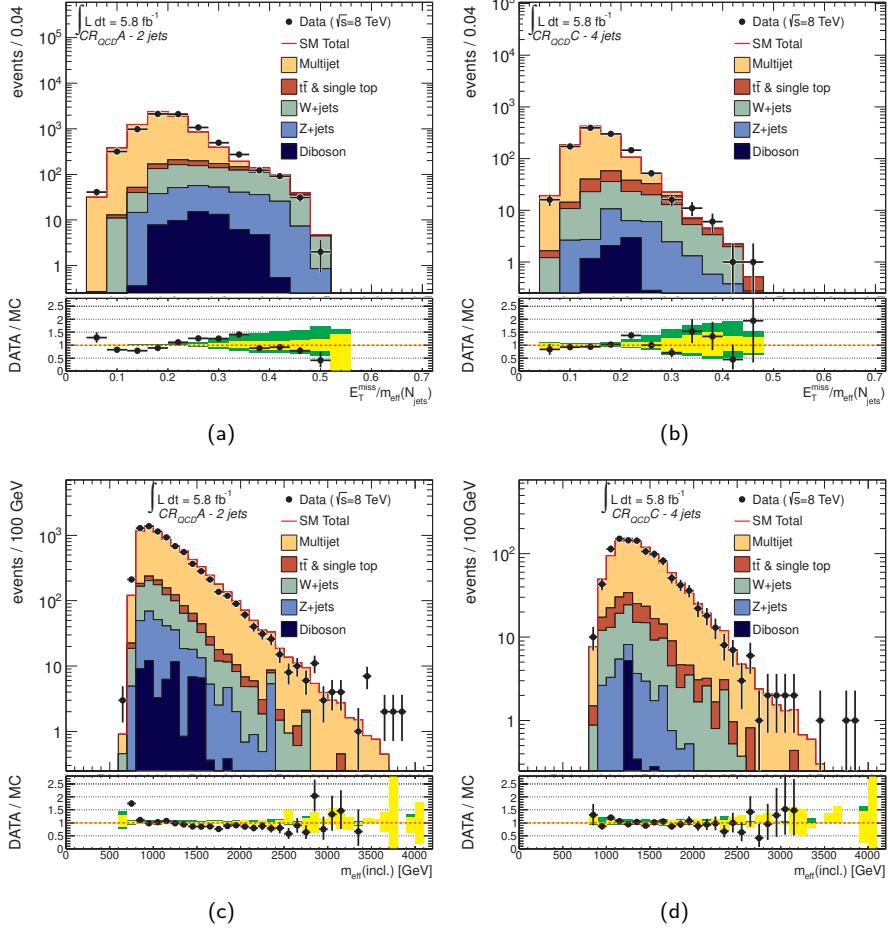


Figure 5.19: Observed and simulated $E_T/m_{\text{eff}}^{\text{excl}}$ and m_{eff} distributions in the 2- and 4-jet channel for CR_{QCD} .

There are two statistical uncertainties on the background which contribute. On the one hand, the limited number of MC events for each sample is a source of uncertainty, given by $\sqrt{N_{\text{MC}}}$. On the other hand, the number of observed events in the control regions propagates into the SR as well, since the observed number of events normalises the background prediction for the corresponding process. A statistical fluctuation thus leads to a fluctuation in the background estimation.

Theoretical uncertainties on the background prediction are numerous, and affect the transfer factors. For $t\bar{t}$ events the main theoretical uncertainties come from uncertainties on the PDF, choice of factorisation and renormalisation scale, matrix element treatment, and parton shower and hadronisation algorithms. The last three are taken into account by comparing the nominal MC generator to an alternative generator.

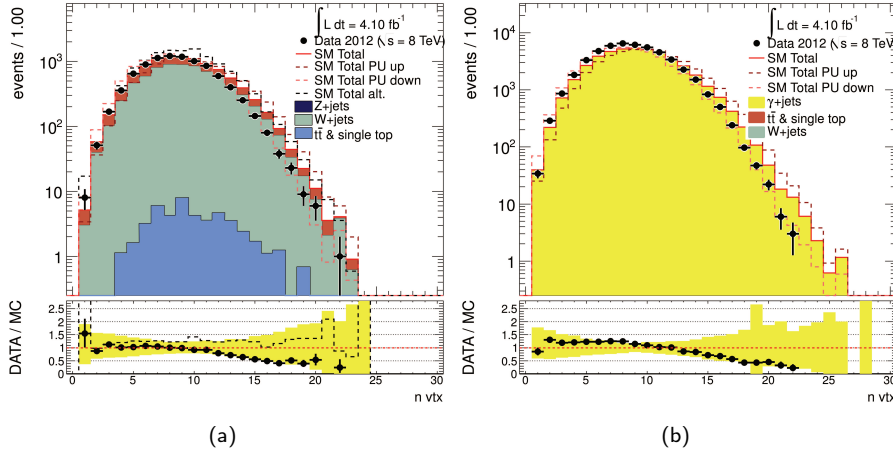


Figure 5.20: Distribution of the number of vertices per event in (a) $CR_W + CR_T$ and (b) CR_γ after requiring $\cancel{E}_T > 160$ GeV and 2 jets. The dashed lines show MC samples when $\langle\mu\rangle$ is scaled with 0.9 and 1.1. The yellow error band in the lower pad, showing the data to MC ratio, includes the experimental uncertainties plus the uncertainty on $\langle\mu\rangle$. Figure taken from [242].

For instance, from the middle left figure 5.15 it is observed that SHERPA $t\bar{t}$ samples (dashed lines) predict too many jets in CR_T while MC@NLO describes the distribution slightly better. The difference between SHERPA and MC@NLO also compares different matching schemes, and different parton shower and hadronisation algorithms. The difference between the prediction of the two generators is therefore taken as a conservative uncertainty on the transfer factor. This leads to uncertainties of the order of 25% in signal regions D and E. As this uncertainty is significantly larger than the uncertainties from PDFs and factorisation/renormalisation scales ($\sim 8\%$), it is taken as the only theoretical uncertainty. For the W +jets and Z +jets MC samples, the theoretical uncertainties come from the choice of factorisation and renormalisation scale, choice of parton shower matching scale, the number of partons generated and the PDF uncertainties. These theoretical uncertainties are evaluated by varying the scales up and down; raising the matching scale; reducing the number of generated partons; and varying the PDF in the generator. After each modification, the transfer factor is recalculated and the quadratic sum of these variations is taken as its uncertainty. For the diboson samples, a flat conservative 50% uncertainty is applied to the WW , WZ and ZZ NLO cross sections obtained with MCFM [255].

For the QCD estimation, a disagreement is seen between data and pseudo-data in the $\Delta\phi(j, \cancel{E}_T)$ distribution before any m_{eff} cut is applied, shown in figure 5.7 (a). Although the disagreement decreases for high m_{eff} values, a conservative uncertainty of 100% is applied to the QCD estimation. This is motivated by the fact that the estimated number of QCD multi-jet events in our SR is very small – in section 5.6 we

Model	$\tilde{q}\tilde{q}$ direct decay		$\tilde{g}\tilde{g}$ direct decay	
	$m_{\tilde{q}} = 450 \text{ GeV},$	$m_{\tilde{q}} = 800 \text{ GeV},$	$m_{\tilde{q}} = 700 \text{ GeV},$	$m_{\tilde{q}} = 1250 \text{ GeV},$
	$m_{\tilde{\chi}_1^0} = 400 \text{ GeV}$	$m_{\tilde{\chi}_1^0} = 50 \text{ GeV}$	$m_{\tilde{\chi}_1^0} = 5500 \text{ GeV}$	$m_{\tilde{\chi}_1^0} = 50 \text{ GeV}$
JES	+12%	+41%	+18%	+8%
JRES	-39%	-37%	-8%	-9%
$\cancel{E}_T$ soft term resolution	$\pm 6\%$	$\pm 3\%$	$\pm 2\%$	$\pm 1\%$
$\cancel{E}_T$ soft term scale	+4%	+0%	+0%	+1%
Pile-up	-4%	-1%	-2%	-0%
	+4%	0%	0%	0%
	-0%	+5%	+6%	+4%
	+6%	-1%	-2%	-2%
	-3%			

Table 5.5: Breakdown of dominant experimental uncertainties on signal samples for 4 points, each in the SR with the highest sensitivity for this point.

will see it is less than 1 event. Thus, such a conservative uncertainty does not affect the fit much.

Finally, the uncertainty on the luminosity measurement is 3.6% as estimated with dedicated ATLAS runs. Uncertainties on the energy scale and resolution of the leptons and photon in the CRs are ignored.

The quadratic sum of the experimental uncertainties is shown in figures 5.15-5.19 per bin as a yellow band in the lower pad, while the green band gives the total uncertainty including the theoretical and statistical uncertainties.

Furthermore, since the b -jet multiplicity is not well reproduced in the SHERPA samples, as explained in section 5.4.4, the overall normalisation of the W +jets sample is uncertain. A conservative uncertainty is placed on these events covering precisely the applied correction factor. This leads to an uncertainty of $0.95/0.89 \approx 7\%$ on W +jets events in CR_W and $(1 - 0.95)/(1 - 0.89) \approx 45\%$ in CR_T .

5.4.7 Uncertainties on signal samples

Signal samples are subject to similar uncertainties as the SM backgrounds. Theoretical uncertainties due to the choice of PDF, strong coupling constant and factorisation and renormalisation scale have been discussed chapter 3, and are shown in figure 5.5. The total theoretical uncertainty from these sources increases from 10% to 50% for increasing squark or gluino mass.

Uncertainties from detector understanding are treated in the same manner as for background samples for each point individually, where again the jet energy scale and resolution, soft terms of the $\cancel{E}_T$ and pile-up uncertainties are the dominant sources. A breakdown of experimental uncertainties is given in table 5.5 for two points in the squark-pair and two in the gluino-pair production simplified model, with for both models one high mass and one compressed point.

However, the theoretical uncertainty on the modelling of ISR jets in the signal MC samples is performed somewhat differently than for background samples. As it contributes significantly to the total signal uncertainty, it is discussed here separately.

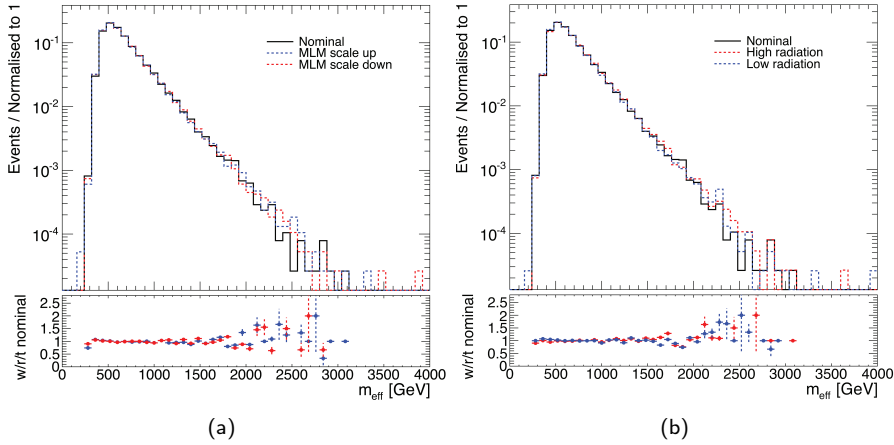


Figure 5.21: Effect of varying (a) the MLM matching scale by a factor of two and (b) the PYTHIA parton shower tune on the m_{eff} distribution for a squark-pair production signal sample with $m_{\tilde{q}} = 500$ GeV and $m_{\tilde{\chi}_1^0} = 350$ GeV in events with no lepton, at least 2 jets and $\cancel{E}_T > 160$ GeV. Both are shown with respect to the nominal sample (in black).

ISR uncertainties

For simplified models with small mass differences between the coloured SUSY particles and the LSP, less momentum is available for the final state quarks coming from the squark decay, and thus it is less likely that one of the quarks forms a jet which passes the trigger requirements. For these compressed SUSY scenarios additional ISR jets are therefore important for the acceptance of SUSY events. Furthermore, a pair of heavy SUSY particles recoiling against an ISR jet are boosted in the same direction. This decreases the angle between them, leading to a final state with higher $\cancel{E}_T$ and $\sum p_T$.

The systematic uncertainties on the modelling of ISR jets in our simplified models have been studied in the ATLAS search for SUSY with one or more leptons at $\sqrt{s} = 7$ TeV [167]. To evaluate the uncertainty from ISR on signal, events are simulated using the matrix element generator MADGRAPH interfaced with PYTHIA for parton showering. For each signal sample the systematic uncertainties are estimated by varying the amount of additional radiation in the parton showers. This is achieved by varying between the nominal PYTHIA sample and two samples with a different ‘tune’: tune 351, which gives high radiation, and tune 352 giving low radiation. In addition, the MLM matching scale used in MADGRAPH to match to the parton shower is varied up and down by a factor 2. Figure 5.21 shows the two variations in our m_{eff} distribution for a sample from the squark-pair production grid with $m_{\tilde{q}} = 500$ GeV and $m_{\tilde{\chi}_1^0} = 350$ GeV. The quadratic sum of the maximum difference in signal acceptance from the two variations (parton shower tunes and MLM matching) is taken as uncertainty [167, 256].

For most points in the parameter space 30000 events have been generated at truth level, while for compressed points with very low selection efficiency due to the small mass difference Δm , the number of generated events is 300000. Yet even with these large number of initial events, only of the order of 100 events survive the signal region selections for the compressed points, and thus the impact of the statistical fluctuations on the ISR uncertainty is large. Fitting the resulting uncertainties on the two dimensional grid of points, a parametrisation has been constructed as function of the gluino, squark and LSP mass:

$$\Delta_{ISR} = 0.48 - 0.0007 \times \min(m_{\tilde{g}}, m_{\tilde{q}}) - 0.0003 \times \Delta m. \quad (5.8)$$

The error increases for lower masses of the squarks and gluinos, and also for smaller mass differences Δm , from negligible for $\min(m_{\tilde{g}}, m_{\tilde{q}}) > 600$ GeV and high mass splitting, to about 30% for low mass splittings and low squark/gluino masses. Although the above method was defined for a leptonic SUSY analysis, it has been confirmed to be applicable for our analysis as well: for 10 points per signal grid the ISR uncertainty was evaluated by varying the PYTHIA tunes and MLM matching scale as described above, which each agree with the parametrisation.

5.5 Statistical procedure for discovery or exclusion

After the measurements and comparisons with theoretical predictions using simulated MC samples in the CRs, the results need to be evaluated. First, the number of estimated background events in each of the SRs is obtained from the CRs and transfer factors. Subsequently, the observed number of events in these SRs are compared with the estimation. Any deviation is interpreted in terms of the probability of there being an excess over the background, which could be a sign of signal events in the SRs. The used statistical procedure takes as input the number of events in the various regions, and gives as output either a discovery p -value, discussed shortly, or an exclusion limit.

To obtain a unified background estimation a likelihood method is used. It combines the measurements in the four CRs to determine the background components ($t\bar{t}$, W +jets, Z +jets and QCD multi-jets) coherently across all regions. In other words, the likelihood fit ensures that the normalisation of the background processes is consistent between all CRs and SR.

With a SR background estimation, including all uncertainties, two hypotheses are tested: the background-only or signal plus background. To quantify differences between observed event counts and the expectation, naively one would use the event counts themselves. Yet the interplay between CRs and SR leads to a complication: the interconnected measurements lead to SR and CR measurements which are not independent. Therefore, deviations from the prediction are not easily quantifiable using the event counts only. To solve this, a variable is introduced which takes into account the correlation between the CRs and the SR. The usual choice for this *test statistic* is the profile log likelihood ratio (LLR), which, as the name suggests, is the ratio of two likelihood functions, each giving the likelihood of the observation given

the estimations and systematic uncertainties, while having a slightly different input as is discussed in section 5.5.3. These likelihoods are each maximised to obtain the best fit to data using three different fits:

'background' fit: To gain confidence in the normalisation procedure, a 'background' fit is performed without information from the SR. In this fit, processes which act as a background in a specific CR are normalised using the data in their own designated CRs. The estimations are also extrapolated to the validation regions, which are used to validate how well the data is described by the estimations from the fit. Note that this fit is not used to estimate the backgrounds in the SR.

'discovery' fit: To search for an excess in any of the SRs, the SR observations and expectations are added to the fit. However, no specific SUSY signal is assumed in either SR or CRs: the 'discovery' fit tests only if more events are observed than expected. It leads to a p -value (p_0) defining the probability that a deviation in data of at least the observed magnitude is due to random fluctuations in the background processes. In addition, a sample with precisely one generic non-SM event is artificially introduced in the SR with floating strength μ , leading to an upper limit on any possible non-SM signal.

'exclusion' fit: In absence of a large excess, certain SUSY signal models are excluded using an 'exclusion' fit. For each SUSY signal point, the SR event count is added to the SR and CRs, and the SR observation is compared to the number of estimated background events plus the signal expectation. The fit gives another p -value (p_1) describing the probability of the observed data being consistent with the estimated signal plus background events.

These three fits are all implemented in the HISTFITTER package, which is based on the HISTFACTORY [257] tool in the ROOSTATS [258] framework. For more information, the full procedure is documented in ref. [259].

Both the likelihood function and the profile log likelihood ratio are discussed shortly. First, the method to obtain the estimated number of events in any region for any process is described.

5.5.1 Estimated number of background events

In this unified approach, all background processes are determined in CRs, together with their cross-contamination: the number of events of background i in the CR dedicated to background j . The estimated total number of events λ_r in a signal or control region r depends on: the number of MC background events $\vec{b}$ in that region (or pseudo events in case of QCD), and the number of raw MC signal events s scaled to the luminosity; the transfer factors $TF_{p,r} = N_{p,r}^{exp} / N_{r',p}^{exp}$ which transfer the number of events for process p from any other region r' to region r ; the signal strength, or signal normalisation factor, μ ; and finally nuisance parameters $\vec{\theta}$ which parametrise the systematic uncertainties. Note that μ is a relative strength: $\mu = 0$ indicates

a background only hypothesis, while $\mu = 1$ is the nominal signal plus background hypothesis. The transfer factors are treated as free parameters in the fit, and scale the measurements of background p in the CR to background estimates in either the signal region, or the cross-contamination in other CRs. This can be written as:

$$\lambda_r(\mu, s, \vec{b}, \vec{\theta}) = \mu \cdot s \cdot TF_{s,r}(\vec{\theta}) + \sum_{p \in \text{backgrounds}} b_p \cdot TF_{p,r}(\vec{\theta}), \quad (5.9)$$

where for the backgrounds without CR (single top and diboson production) the scale factor and the transfer factors are fixed to 1.

In our analysis, there are 4 CRs for each SR, which each target one process. For these processes, the transfer factors are not fixed in the fit to the exact ratio between the regions as given by equation 5.4, but are left to move slightly to obtain the best fit to all background processes simultaneously between the CRs. There are 15 nuisance parameters per transfer factor, while the number of raw background events in each of the CRs and SR are taken as input.

To combine all the regions and make full use of the nuisance parameters, a likelihood function is introduced.

5.5.2 The likelihood function

The likelihood function L gives the likelihood of $\vec{n}$, the set of measurements n_r in the signal region or each of the control and validation regions r , given the number of expected background and signal events $\vec{b}$ and $\vec{s}$ and nuisance parameters θ . It is defined for each of the SRs separately as:

$$L(\vec{n}|\mu, \vec{s}, \vec{b}, \vec{\theta}) = \prod_{r \in [SR, CRs]} P_r(n_r|\lambda_r(\mu, s_r, \vec{b}, \vec{\theta})) \prod_{s \in \text{sys.unc.}} G(\theta_s^0 - \theta_s). \quad (5.10)$$

The Poisson probability density functions (pdfs) $P_r(n_r|\lambda_r(\mu, s_r, \vec{b}, \vec{\theta}))$ define the probability of measuring n_r events in region r assuming λ_r expected events. The second term is a product over Gaussian pdfs which constrain the nuisance parameters $\vec{\theta}$. The nuisance parameters for systematic uncertainty s are written as a nominal value θ_s^0 , around which can be varied by a value θ_s . The nominal value may be taken as 0, which leads to variations θ_s corresponding to ± 1 standard deviations around the nominal value. By maximising the likelihood function in equation 5.10 given measurements $\vec{n}$, the transfer factors, nuisance parameters $\vec{\theta}$ and signal strength μ are fixed to values which are most consistent with $\vec{n}$. After maximisation of the likelihood, the expected number of events of each background process is determined in each of the CRs and SR.

5.5.3 Profile log likelihood ratio

There are two options when interpreting the results, i.e. the comparison of the observed versus the estimated number of events in the SR. When evaluating if a discovery has

been made, the Standard Model hypothesis is tested against any excess above the estimated background events, without assuming any signal in the SRs. Alternatively, when no discovery can be claimed, we try to exclude models. Here the signal hypothesis is tested, assuming a specific signal model and taking thus $\mu = 1$.

As mentioned previously, the usual choice of test statistic is the profile log likelihood ratio, which is defined as:

$$\Lambda(\mu, \vec{n}, \vec{\theta}^0) = -2 \ln \frac{L(\vec{n}, \vec{\theta}^0 | \mu, \hat{\vec{b}}, \hat{\vec{\theta}})}{L(\vec{n}, \vec{\theta}^0 | \hat{\mu}, \hat{\vec{b}}, \hat{\vec{\theta}})}, \quad (5.11)$$

where $\hat{\mu}, \hat{\vec{b}}, \hat{\vec{\theta}}$ are the values of $\mu, \vec{b}$ and $\vec{\theta}$ which maximise L for the observed $\vec{n}$, and the double-hatted $\hat{\vec{b}}, \hat{\vec{\theta}}$ maximise L for the specific value of μ , given data $\vec{n}$.

We assume that a SUSY signal leads to an increase of SR events, while a decrease is deemed unphysical. Therefore, we set the requirement on the test statistic q_0 when searching for a discovery to be equal to zero when the number of events is below the background estimation [260]:

$$q_0 = \begin{cases} -2 \ln \Lambda(0, \vec{n}, \vec{\theta}^0) & \text{if } \hat{\mu} \geq 0, \\ 0 & \text{if } \hat{\mu} < 0, \end{cases} \quad (5.12)$$

For exclusion, where the signal strength $\mu = 1$, models should not be excluded when more events are observed than expected from the signal plus background hypothesis, thus when the best-fit value of the signal strength is larger than 1. We can therefore define the test statistic q_1 as

$$q_1 = \begin{cases} -2 \ln \Lambda(1, \vec{n}, \vec{\theta}^0) & \text{if } \hat{\mu} \geq 1, \\ 0 & \text{if } \hat{\mu} < 1. \end{cases} \quad (5.13)$$

Given the above test statistics, the p -value, giving the probability of obtaining the observed deviation from the null-hypothesis by random fluctuations, is defined as

$$p_\mu = \int_{q_{\mu, \text{obs}}}^{\infty} f(q_\mu | \mu) dq_\mu, \quad (5.14)$$

with $q_{\mu, \text{obs}}$ the observed value of the test statistic and $f(q_\mu | \mu)$ the pdf of test statistic q_μ given signal strength μ . The distribution $f(q_\mu | \mu)$ is not easily calculated exactly, but can be approximated by an asymptotic formula using the Asimov method [260], or by generating pseudo-data from the likelihood using toy Monte Carlo techniques. For discovery tests, we compare with the background only hypothesis where $\mu = 0$, and thus look at p_0 , while for exclusion, the signal plus background hypothesis is tested using p_1 .

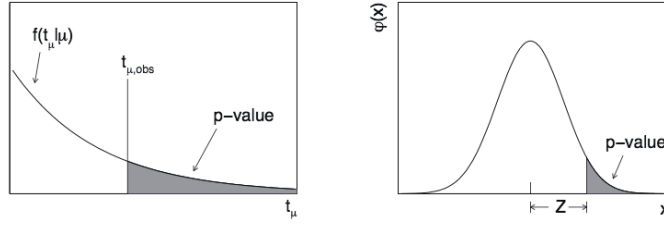


Figure 5.22: Illustration of (a) the definition of the p -value, given a pdf f of a test statistic t_μ and an observed value t_μ^{obs} and (b) the standard normal distribution ϕ and the definition of the significance Z and the p -value. Figure taken from [260].

Significance

Having calculated the p -value for each of the SRs, it is convenient to work with the *significance* Z_0 of the result, given by

$$Z_0 = \Phi^{-1}(1 - p_0), \quad (5.15)$$

with Φ the cumulative distribution function of the standard normal distribution, and its inverse Φ^{-1} is the quantile function. See figure 5.22 for an illustration of the relationship. It is common to use the rule that $Z_0 \geq 3$ defines *evidence* for a discovery, while $Z_0 \geq 5$ or $p_0 \leq 2.87 \times 10^{-7}$ is declared to be a *discovery*. On the other hand, if no large excess is observed, p_1 is very small. It would be straightforward to consider a model as excluded at 95% confidence level if $p_1 \leq 0.05$. Yet to protect against models being excluded to which the analysis has no sensitivity, for example by an under-fluctuation of the background, the CL_s method is introduced [261]. By defining an exclusion only if

$$\text{CL}_s \equiv \frac{p_1}{1 - p_0} \leq 0.05, \quad (5.16)$$

models with a very small signal expectation but also little difference between signal and background will not be excluded. It should be noted that CL_s is not a real probability, and it *overcovers*⁴ and thus is conservative. CL_s is therefore only used for setting exclusion limits, while p_0 is used for discovery tests.

5.6 Results

The results of the analysis are obtained by gathering all the information of the previous sections together. Since none of the validation regions has a significant disagreement between observation and prediction, we can look at the SRs. The observed m_{eff}

⁴The coverage of a confidence interval is the fraction of repetitions of the experiment in which the interval contains the true value.

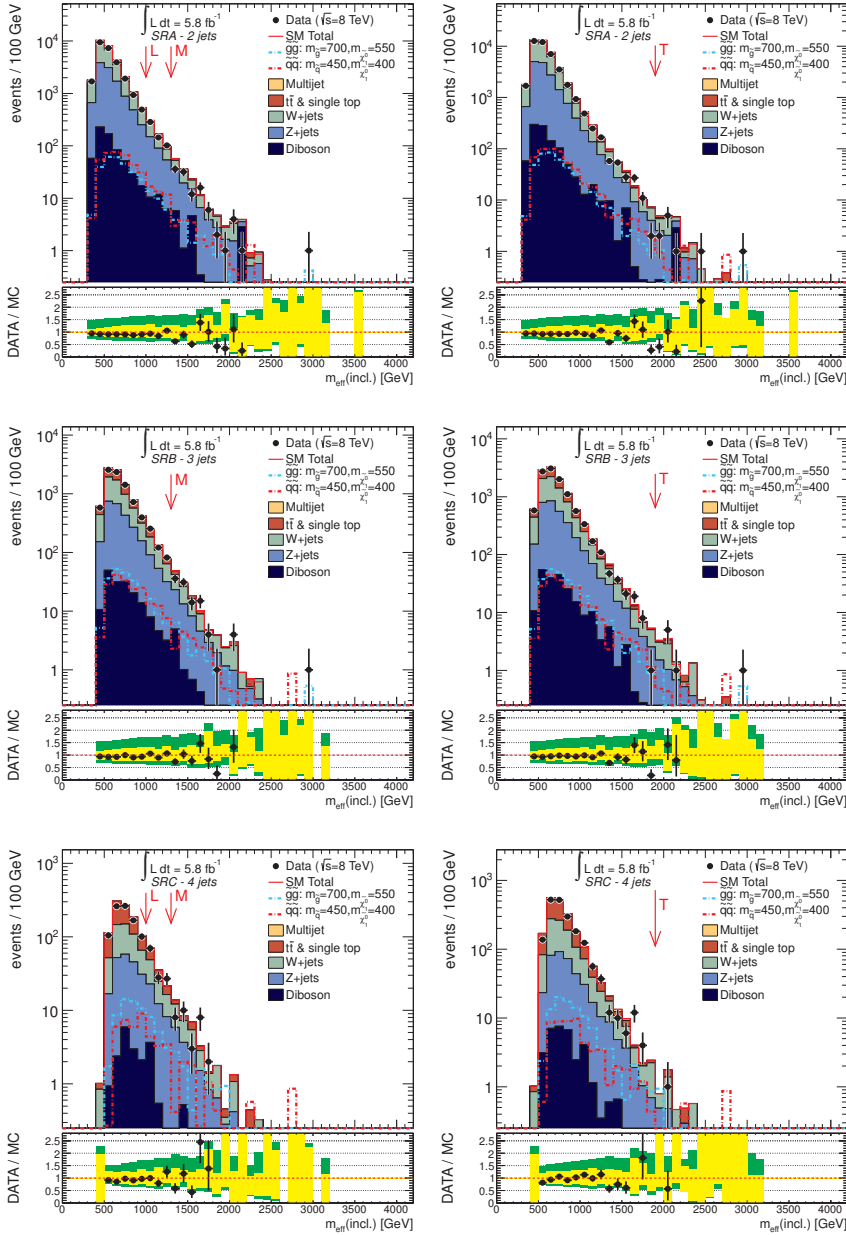


Figure 5.23: Observed final m_{eff} distributions, for SRA 'loose' and 'medium' (top left), SRA 'tight' (top right), SRB 'medium' (middle left), SRB 'tight' (middle right), SRC 'loose' and 'medium' (bottom left) and SRC 'tight' (bottom right). The red arrows indicate where the final cut on m_{eff} is placed for each of the SRs. The dashed histograms indicate a compressed $\tilde{g}\tilde{g}$ (light blue) and a $\tilde{q}\tilde{q}$ (red) SUSY signal.

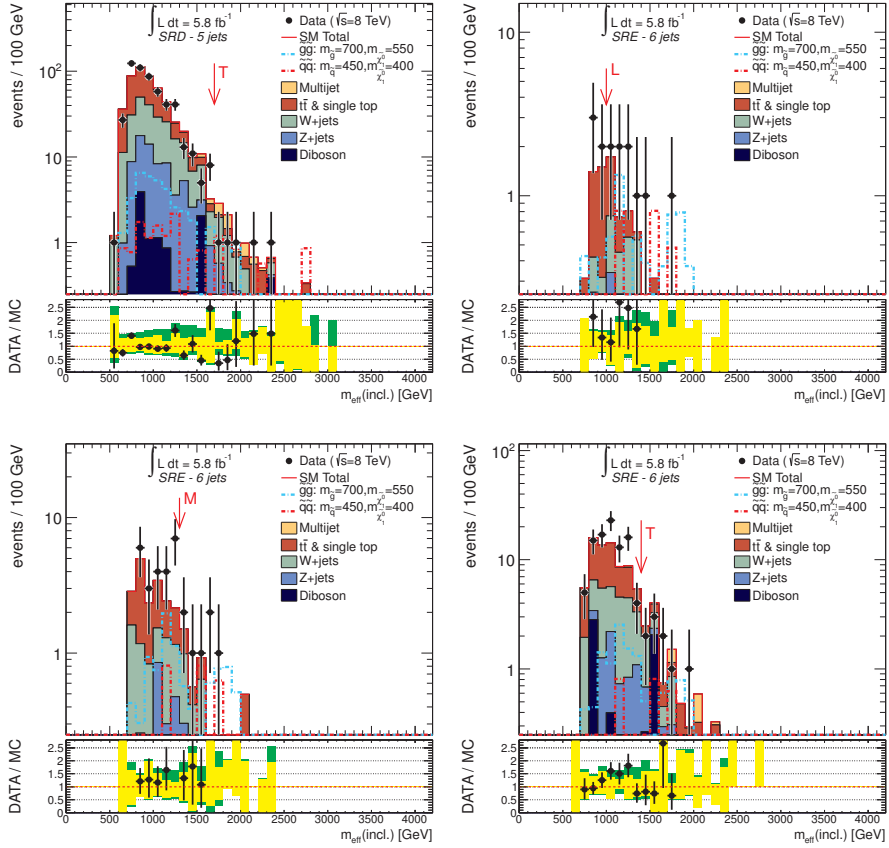


Figure 5.24: Observed m_{eff} distributions after all cuts, for SRD ‘tight’ (top left), SRE ‘loose’ (top right), SRE ‘medium’ (bottom left) and SRE ‘tight’ (bottom right). The red arrows indicate where the final cut on m_{eff} is placed for each of the SRs. The dashed histograms indicate a compressed gluino-pair (light blue) and a squark-pair (red) SUSY signal.

distribution in each of the SRs is shown in figures 5.23 and 5.24 after the $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ cut. Note that the background shown comes from raw MC simulation. For the final result, the likelihood fit as described in section 5.5 is performed for the SM background estimation, comparing its results with the observed data in the SRs. The result for each of the SRs is shown in tables 5.6 and 5.7. For each SR the raw number of MC events for each of the backgrounds is given, where the number of QCD multi-jet background events is missing as these are obtained from pseudo data. In the second block the number of background events obtained from the likelihood fit is given with its corresponding uncertainty. The difference between the raw MC prediction and that

	A loose	A med	B med	C loose	C med	E loose	E med
<i>Raw MC prediction:</i>							
Z +jets	338.2	74.7	50.4	55.0	16.1	1.0	0.8
W +jets	264.1	53.5	51.9	62.9	16.4	2.1	1.9
$t\bar{t}$ + single top	74.9	8.1	14.2	42.6	5.3	2.1	1.6
Diboson	53.1	18.2	11.1	6.2	0.9	0	0
<i>Prediction from fit:</i>							
Z +jets	354 $\pm$ 21	81 $\pm$ 8	59 $\pm$ 6	67 $\pm$ 6	18.5 $\pm$ 3.0	2.0 $\pm$ 1.0	0.6 $\pm$ 0.5
W +jets	180 $\pm$ 140	33 $\pm$ 35	32 $\pm$ 34	40 $\pm$ 40	8 $\pm$ 8	1.2 $\pm$ 1.3	0.9 $\pm$ 1.1
$t\bar{t}$ + single top	67 $\pm$ 16	7.6 $\pm$ 3.5	14 $\pm$ 5	39 $\pm$ 7	5.3 $\pm$ 2.0	2.5 $\pm$ 0.9	2.0 $\pm$ 1.4
Diboson	53 $\pm$ 23	18 $\pm$ 9	11 $\pm$ 6	6 $\pm$ 4	0.9 $\pm$ 0.6	–	–
QCD Multi-jets	0.6 $\pm$ 0.6	0.1 $\pm$ 0.1	0.2 $\pm$ 0.2	–	–	–	–
Total background	650 $\pm$ 130	140 $\pm$ 33	115 $\pm$ 30	155 $\pm$ 31	33 $\pm$ 8	5.7 $\pm$ 1.7	3.5 $\pm$ 1.7
Observed	643	111	106	156	31	9	7
Discovery p -value	0.498	0.500	0.500	0.486	0.498	0.161	0.108
$\sigma_{BSM}^{\text{eff}}$ U.L. [fb]	38.8	5.84	7.55	11.3	3.09	1.79	1.71

Table 5.6: Number of raw MC and fitted background events in each of the ‘loose’ and ‘medium’ SRs, compared to the observed number of events. Empty cells in the background estimations correspond to less than 0.01 estimated events. The last two rows give the discovery p -value p_0 , and the 95% CL upper limit on the cross section of any beyond-the-Standard Model process still allowed in fb.

from the fit comes from the transfer factors between the CRs and the SRs. Finally, in the last rows of tables 5.6 and 5.7 the observed number of events in each signal region are given, together with the p -value p_0 which is obtained from the ‘discovery’ fit. A breakdown of the systematic uncertainties on the total number of SM background events is shown in table 5.8. The statistical uncertainty in each of the SRs is calculated as $\sqrt{N_{\text{obs}}}$.

None of the SRs shows any significant excess of the observed number of events over the SM estimation, as is seen from the p -values which are mostly very close to 0.5. Only SRE ‘loose’ and ‘medium’ show a deviation from the expected number of events, with $p_0 = 0.161$ and $p_0 = 0.108$, respectively, which correspond to a significance of $Z_0 = 0.99$ and $Z_0 = 1.24$. Before interpreting these results, several observations are made from the tables.

First of all, the SHERPA heavy flavour uncertainty on W +jets applied in CR_W and CR_T is the dominant uncertainty in the ‘loose’ and ‘medium’ selections, especially in

	A tight	B tight	C tight	D tight	E tight
<i>Raw MC prediction:</i>					
Z +jets	7.4	4.5	1.9	1.3	1.3
W +jets	6.6	5.6	2.1	3.4	3.3
$t\bar{t}$ + single top	1.0	1.1	0.6	1.8	2.7
Diboson	3.3	0.2	0.0	0.8	2.6
<i>Prediction from fit:</i>					
Z +jets	6.8 ± 2.2	5.1 ± 1.7	2.0 ± 1.1	2.5 ± 1.1	1.2 ± 0.7
W +jets	3 ± 4	2.7 ± 3.4	0.3 ± 0.5	–	0.8 ± 1.3
$t\bar{t}$ + single top	0.8 ± 0.8	0.8 ± 0.9	0.6 ± 0.5	2.6 ± 1.6	5.1 ± 3.3
Diboson	3.3 ± 3.1	0.2 ± 1.4	–	0.8 ± 0.4	2.6 ± 2.0
QCD Multi-jets	–	–	–	0.4 ± 0.5	0.1 ± 0.2
Total background	14 ± 5	8.7 ± 3.4	2.8 ± 1.2	6.3 ± 2.1	10 ± 4
Observed	10	7	1	5	9
Discovery p -value	0.499	0.500	0.499	0.500	0.499
$\sigma_{BSM}^{\text{eff}}$ U.L. [fb]	1.53	1.26	0.57	1.03	1.60

Table 5.7: Number of raw MC and fitted background events in each of the ‘tight’ SRs, compared to the observed number of events. Empty cells in the background estimations correspond to less than 0.01 estimated events. The last two rows give the discovery p -value p_0 , and the 95% CL upper limit on the cross section of any beyond-the-Standard Model process still allowed in fb.

SRA, SRB and SRC. As this uncertainty is only applied in the control region, and not in the signal region, it does not cancel in the transfer factor, and thus its relative effect increases with respect to the other systematics which do partially cancel. This will affect mainly the low and intermediate mass scenarios, and compressed regions, as here the ‘loose’ and ‘medium’ selections are most sensitive.

Secondly, there are only between 5 and 10 events observed in the ‘tight’ CR selections, even though in CR_W and CR_T the $\Delta\phi$ and $E_T/m_{\text{eff}}^{\text{excl}}$ requirements have not been applied. Therefore, the yields in these control regions constitute a large contribution to the total uncertainty in the ‘tight’ SRs. This is the downside of having defined CRs as close to the SR selections as possible. The same holds for the MC statistics, which are of the order of 100 events in the ‘tight’ selections of the CRs, and thus contribute significantly to the total uncertainty.

The systematic uncertainties due to the jet energy scale and resolution are relatively small in most ‘loose’ and ‘medium’ SRs (of the order of 10%), while in the ‘tight’

	A loose	A med	B med	C loose	C med	E loose	E med
Total systematic	± 129.0	± 33.1	± 30.3	± 30.8	± 7.5	± 1.7	± 1.7
Total statistical	± 25.4	± 10.5	± 10.3	± 12.5	± 5.6	± 3.0	± 2.6
MC statistics	± 0.0	± 0.0	± 0.0	± 0.0	± 0.0	± 0.56	± 0.37
Factorisation scale	± 0.06	± 0.04	± 0.06	± 0.04	± 0.03	± 0.01	± 0.00
Renormalisation scale	± 0.04	± 0.02	± 0.04	± 0.03	± 0.02	± 0.00	± 0.00
Jet energy resolution	± 0.03	± 0.05	± 0.11	± 0.03	± 0.19	± 0.05	± 0.60
Jet energy scale	± 0.12	± 0.04	± 0.06	± 0.12	± 0.03	± 0.06	± 0.00
Pile-up	± 0.04	± 0.09	± 0.04	± 0.03	± 0.15	± 0.18	± 0.08
Soft $\cancel{E}_T$ resolution	± 0.01	± 0.00	± 0.00	± 0.03	± 0.01	± 0.11	± 0.13
Soft $\cancel{E}_T$ scale	± 0.01	± 0.01	± 0.00	± 0.02	± 0.02	± 0.04	± 0.09
Diboson theory	± 0.13	± 0.24	± 0.16	± 0.09	± 0.06	± 0.01	± 0.00
Top theory	± 0.02	± 0.02	± 0.00	± 0.10	± 0.03	± 0.13	± 0.24
W +jets SHERPA correction	± 0.97	± 0.88	± 0.93	± 0.96	± 0.80	± 0.53	± 0.44
Multi-jets method	± 0.00	± 0.00	± 0.01	± 0.00	± 0.00	± 0.00	± 0.00
Multi-jets CR statistics	± 0.00	± 0.00	± 0.00	± 0.00	± 0.00	± 0.00	± 0.00
W +jets CR statistics	± 0.17	± 0.20	± 0.23	± 0.25	± 0.28	± 0.31	± 0.35
Z +jets CR statistics	± 0.16	± 0.21	± 0.19	± 0.20	± 0.39	± 0.49	± 0.28
$t\bar{t}$ CR statistics	± 0.13	± 0.09	± 0.16	± 0.23	± 0.19	± 0.31	± 0.43
	A tight	B tight	C tight	D tight	E tight		
Total systematic	± 5.2	± 3.4	± 1.2	± 2.1	± 3.7		
Total statistical	± 3.2	± 2.6	± 1.0	± 2.2	± 3.0		
MC statistics	± 0.42	± 0.22	± 0.31	± 0.40	± 0.61		
Factorisation scale	± 0.02	± 0.04	± 0.02	± 0.01	± 0.00		
Renormalisation scale	± 0.01	± 0.02	± 0.00	± 0.00	± 0.00		
Jet energy resolution	± 0.08	± 0.08	± 0.18	± 0.50	± 0.36		
Jet energy scale	± 0.12	± 0.22	± 0.17	± 0.31	± 0.15		
Pile-up	± 0.46	± 0.06	± 0.13	± 0.01	± 0.06		
Soft $\cancel{E}_T$ resolution	± 0.02	± 0.03	± 0.03	± 0.00	± 0.02		
Soft $\cancel{E}_T$ scale	± 0.01	± 0.02	± 0.01	± 0.07	± 0.08		
Diboson theory	± 0.28	± 0.02	± 0.00	± 0.18	± 0.36		
Top theory	± 0.04	± 0.05	± 0.03	± 0.12	± 0.26		
W +jets SHERPA correction	± 0.50	± 0.67	± 0.19	± 0.00	± 0.20		
Multi-jets method	± 0.00	± 0.01	± 0.00	± 0.21	± 0.04		
Multi-jets CR statistics	± 0.00	± 0.00	± 0.00	± 0.03	± 0.02		
W +jets CR statistics	± 0.35	± 0.49	± 0.31	± 0.00	± 0.27		
Z +jets CR statistics	± 0.34	± 0.47	± 0.87	± 0.45	± 0.18		
$t\bar{t}$ CR statistics	± 0.15	± 0.25	± 0.30	± 0.47	± 0.54		

Table 5.8: Breakdown of the dominant uncertainties on the number of fitted background events in each of the SRs. The individual uncertainties are given relative to the total uncertainty. Note that as the uncertainties can be correlated, the quadratic sum does not need to add up to 1.

selections (except for SRA) they range between 10 and 50%. Yet the uncertainty on JES and JRES are taken from the variation of these quantities in the MC samples, and thus depend heavily on the MC statistics. Since the MC statistics become small in the ‘tight’ regions, the JES and JRES uncertainties increase as well.

The impact of theoretical uncertainties is small for the factorisation and renormalisation scale, and theory uncertainties on W +jets and Z +jets samples. The $t\bar{t}$ theory uncertainty, assessed by comparing MC@NLO and SHERPA samples, is small ($< 10\%$) for SRs up to 5 jets, while it is $\sim 25\%$ for SRE. The only large theoretical uncertainty is for the diboson processes, coming from the conservative 50% uncertainty applied to the SHERPA diboson samples. As no CR was defined for this background, it is included as a pure MC estimate, and is not profiled via the fit, resulting in a remaining large uncertainty. Furthermore, diboson production contributes as much to the total background prediction as $t\bar{t}$ production in the low jet multiplicity channels, at around 10%. The large theoretical uncertainty therefore leads to diboson production having the largest absolute uncertainty after W +jets.

Finally, although an uncertainty of 100% has been applied to the QCD multi-jets estimation, the final uncertainty on the total number of SM background events due to this method is negligible in most SRs, except for SRD.

5.6.1 Interpretation of the results

As no excess has been observed, limits are set on any possible SUSY model still allowed. From the ‘discovery’ fit, which assumes no signal contamination in the CRs, a model independent upper limit on the *effective cross section* $\sigma_{BSM}^{\text{eff}}$ of any beyond-the-Standard Model (BSM) process is obtained, which is given for each SR in tables 5.6 and 5.7. These limits correspond to the maximum allowed cross section for a process on top of the SM background estimate, folded with the acceptance (A) and efficiency (ϵ) of the BSM signal: $\sigma_{BSM}^{\text{eff}} = \sigma_{BSM} \times A \times \epsilon$. They are used by theoreticians to evaluate how their own SUSY model holds against this analysis, given that they have a cross section for their model, together with its kinematic acceptance and selection efficiency.

To interpret the results as an exclusion limit on the SUSY models, an ‘exclusion’ fit is performed, which compares the number of expected signal events with the estimated SM background, and calculates the CL_s for each SUSY signal sample independently. Signal contamination in the CRs is taken into account in this fit. The signal events are taken from MC, with the same experimental sources of systematic uncertainties as the background. In addition, a conservative uncertainty on the production of ISR jets is applied, as described in section 5.4.6. Unlike for SM background, the theoretical uncertainties on the signal production cross section are not treated as nuisance parameters in the likelihood fit – rather, three observed exclusion limits are calculated, corresponding to the nominal value and the ± 1 standard deviation variations of the signal cross section. The combination of the limits in the various SRs is performed by taking for each signal model point the best expected SR, i.e. the SR with the lowest expected CL_s . This results in the optimal analysis, while not biasing the observed limit by cherry picking those regions where less events are observed than expected.

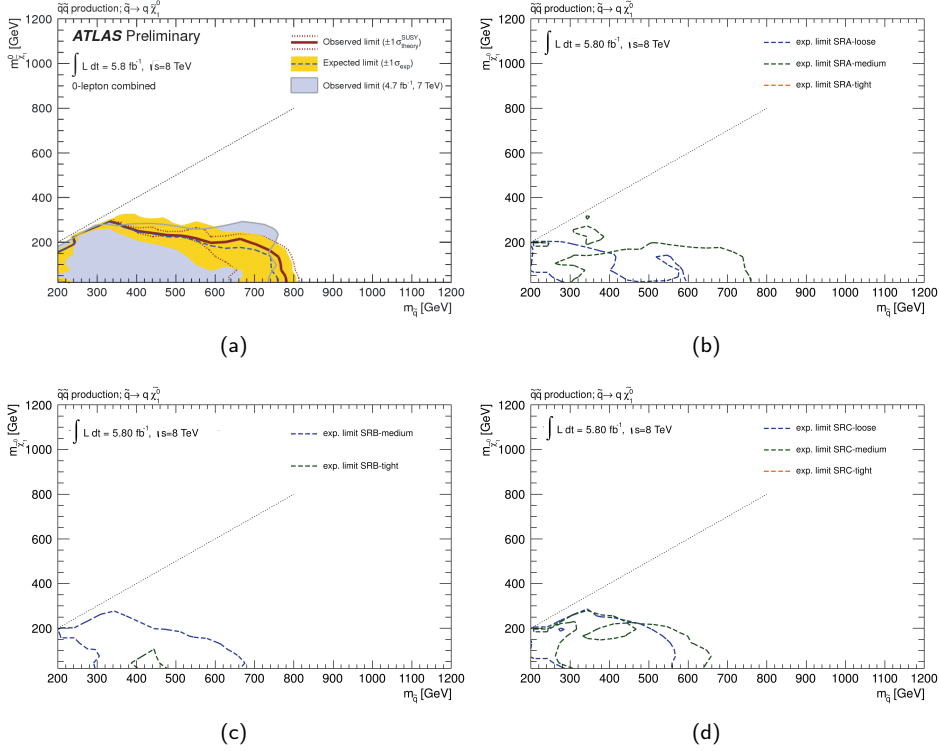


Figure 5.25: (a) The observed 95% CL exclusion limit with 5.8 fb^{-1} of $\sqrt{s} = 8$ TeV data for the simplified model with a directly decaying squark-pair. The expected exclusion limit is indicated by the blue dashed line, with its experimental uncertainties given by the yellow band. The observed limit is indicated by the solid maroon line, while the dotted lines represent the limit when the signal expectation is deviated by its $\pm 1 \sigma$ theoretical uncertainty. The purple shaded area shows the observed limit from the $\sqrt{s} = 7$ TeV analysis. Figures (b)-(d) show the expected exclusion limits for the 2-, 3- and 4-jet SRs.

The combined 95% CL exclusion limits on the simplified models with a direct decay to quarks and LSPs are shown in figures 5.25-5.27, together with the limit of each of the individual SRs in these models. The 5- and 6-jet SRs are not shown for the disquark grid since they are not sensitive to this model. Figure 5.28 shows the best expected SRs for each signal point in the three grids. Figure 5.29 shows the limit on an CMSSM model with $\tan \beta = 10$, $A_0 = 0$ and $\mu > 0$ in the $(m_0, m_{1/2})$ plane and in the $(m_{\tilde{g}}, m_{\tilde{q}})$ plane. As the uncertainties on the estimated number of background events are large, a very large excess or deficit is needed to deviate from the expected exclusion. Since the observed number of events only exceeds the SM prediction in SRs E 'loose' and 'medium', which are not among the best expected SRs, the observed

limit lies close to the expected limit for these models.

The observed mass limit on light squarks in the squark-pair production scenario with decoupled gluinos (figure 5.25 (a)) is set to $m_{\tilde{q}} > 780$ GeV for a massless neutralino, while the LSP mass limit is $m_{\tilde{\chi}_1^0} \gtrsim 300$ GeV for 340 GeV squarks, decreasing to $m_{\tilde{\chi}_1^0} \gtrsim 150$ GeV for more massive squarks, with mass up to 760 GeV. The limit comes very close to the diagonal, where squarks and LSP have equal mass. As no signal samples were produced with $\Delta m(\tilde{q}, \tilde{\chi}_1^0) < 25$ GeV, no limit has been set for smaller mass differences. This limit is similar to the 2011 $\sqrt{s} = 7$ TeV exclusion limit, which is indicated by the purple shaded area. It should be noted that this is somewhat deceptive: the observed limit in the 2011 analysis was significantly stronger than the expected limit due to having seen less events than expected in most SRs. When comparing the expected exclusion limits, our analysis increases the limit on squarks by ~ 50 GeV for low mass LSPs. For higher LSP masses, and thus lower mass splittings, the 2011 limit is ~ 25 GeV higher in the LSP mass between a squark mass of 400 and 550 GeV. The fact that the $\sqrt{s} = 8$ TeV analysis was not able to improve over the $\sqrt{s} = 7$ TeV analysis is due to the increased SM background predictions and their uncertainties in the low jet multiplicity channels, most notably the heavy-flavour bug in SHERPA, while these low jet signal regions are the most sensitive to these models. Furthermore, the uncertainty on the production of additional ISR jets for signal events becomes high for the more compressed region, decreasing the sensitivity. Since the size of the recorded dataset increased only with $\sim 20\%$, higher background uncertainties and similar signal ISR uncertainties lead to similar results to 2011. We conclude that the high systematic error limits the sensitivity. For the compressed region, near the diagonal, the limit is driven by SRC ‘loose’ and SRB and SRC ‘medium’, as is seen from figures 5.25 (b)-(d). The high mass region is driven by SRA and SRB ‘medium’. This is also illustrated by the SRs with the best expected CL_s as seen in figure (e).

In the simplified model with squark-gluino production (figure 5.26), where the light squarks are set to $m_{\tilde{q}} = 0.96 \times m_{\tilde{g}}$, the mass limit on the gluino is $m_{\tilde{g}} > 1440$ GeV for massless neutralinos, while the limit on the neutralino mass decreases from ~ 700 GeV to ~ 550 GeV for $m_{\tilde{g}} < 1400$ GeV. As seen in the squark-pair model, the limit approaches the diagonal of equal mass gluino and LSP, with only a small gap from $m_{\tilde{g}} > 600$ GeV onwards. The limit in the very compressed region is driven by SRA ‘medium’, as seen from the best expected SRs in figure 5.28 (b), where along the total diagonal this is the best performing SR. For intermediate masses ($800 \text{ GeV} < m_{\tilde{q}} < 1200 \text{ GeV}$) the limit is driven by SRC ‘medium’ and for high masses by the tight selections of SRA and SRB, which is also illustrated in figures 5.26 (b)-(d).

The limit on pure gluino-pair production with decoupled squarks shows an improvement over the $\sqrt{s} = 7$ TeV analysis in the high mass range: for massless LSPs, the mass limit is set to $m_{\tilde{g}} \gtrsim 1175$ GeV, an improvement of more than 200 GeV, as is seen in figure 5.27. In the compressed region, gluino-pair production is excluded for models with LSPs up to $m_{\tilde{\chi}_1^0} \gtrsim 500$ GeV, while this limit degrades to 375 GeV for gluinos

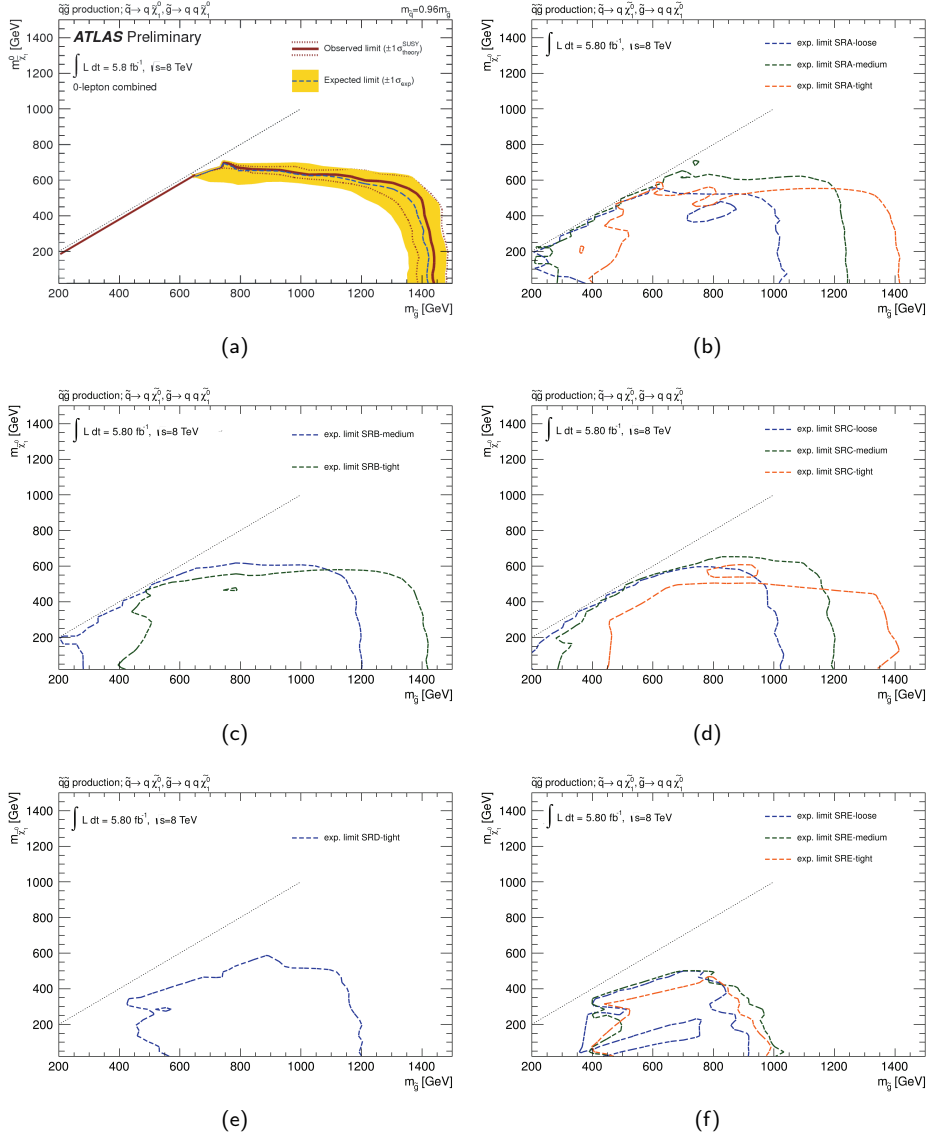


Figure 5.26: (a) The observed 95% CL exclusion limit with 5.8 fb^{-1} of $\sqrt{s} = 8 \text{ TeV}$ data, for the simplified model with squark-gluino production, decaying directly to quarks and LSPs. The expected exclusion limit is indicated by the blue dashed line, with its uncertainties given by the yellow band. The observed limit is indicated by the solid maroon line, while the dotted lines represent the limit when the signal expectation is deviated by its $\pm 1 \sigma$ theoretical uncertainty. Figures (b)-(e) show the expected exclusion limits for the various channels.

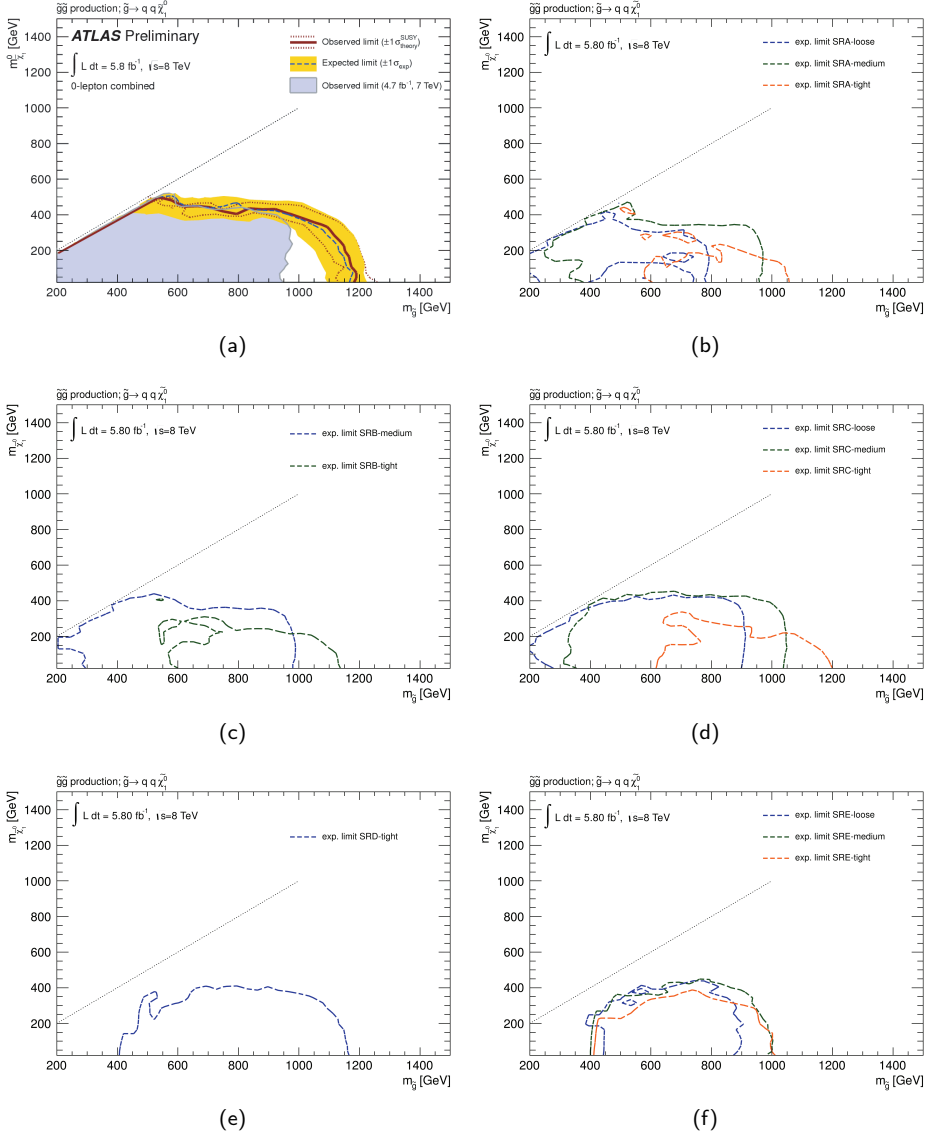


Figure 5.27: (a) The observed 95% CL exclusion limit with 5.8 fb^{-1} of $\sqrt{s} = 8$ TeV data, for the simplified model with gluino-pair production, decaying directly to quarks and LSPs. The expected exclusion limit is indicated by the blue dashed line, with its uncertainties given by the yellow band. The observed limit is indicated by the solid maroon line, while the dotted lines represent the limit when the signal expectation is deviated by its $\pm 1\sigma$ theoretical uncertainty. The purple shaded area shows the observed limit from the 2011 $\sqrt{s} = 7$ TeV analysis. Figures (b)-(e) show the expected exclusion limits for the various channels.

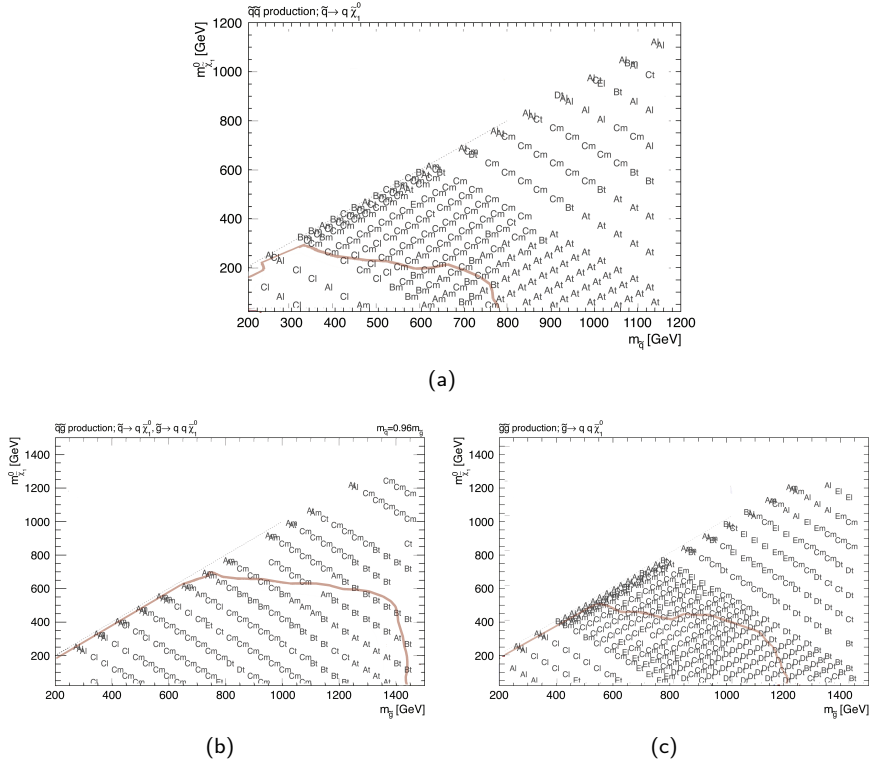


Figure 5.28: The best expected SRs per signal point for the three simplified direct decay models, with (a) squark-pair, (b) squark-gluino and (c) gluino-pair production. The observed limit is shown in the background for reference.

with higher masses. This observed limit lies on top of the excluded region from the 4.7 fb^{-1} analysis on $\sqrt{s} = 7 \text{ TeV}$ data. The best performing SR for the compressed region is again SRA ‘medium’, while for somewhat larger gluino masses the ‘loose’ and ‘medium’ SRC are most sensitive (see figures 5.27 (c) and 5.28 (c)). In the high gluino mass region, SRC and SRD ‘tight’ are most sensitive, with the 5-jet channel (SRD) having the lowest expected CL_s for most points.

Finally, the exclusion limit in the CMSSM plane has improved significantly over the whole plane. The phenomenology of the CMSSM allows for all SUSY processes to occur, with mostly squark-pair production in the low m_0 region, and squark-gluino and gluino-pair production for higher m_0 values, with possibilities of cascade decays. With a higher probability of many jets, SRE ‘tight’ is the best expected SR for high m_0 values, while SRD ‘tight’ is the best expected selection for intermediate and SRC ‘tight’ for low m_0 values, as seen in figure 5.29 (b). In figure 5.29 (c) the mass limits on squarks and gluinos are shown: when degenerate, the mass limit is at $m_{\tilde{g}} = m_{\tilde{q}} = 1500 \text{ GeV}$. This is higher than the limits obtained in the simplified models, as here

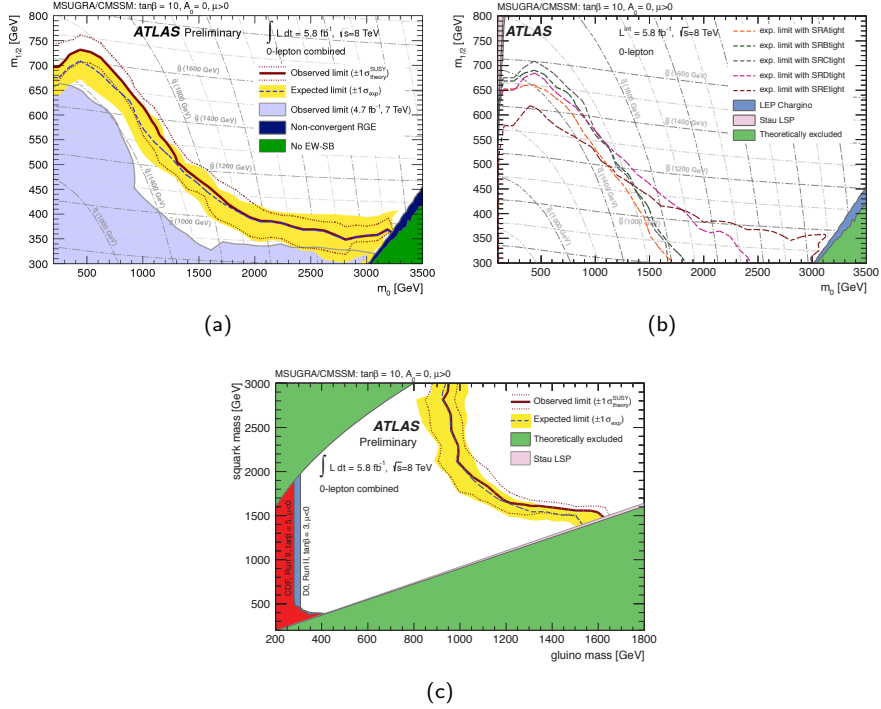


Figure 5.29: The observed 95% CL exclusion limit with 5.8 fb^{-1} of $\sqrt{s} = 8 \text{ TeV}$ data for CMSSM models with $\tan\beta = 10$, $A_0 = 0$ and $\mu > 0$ presented in (a) the m_0 - $m_{1/2}$ plane, and (b) in the m_0 - m_0 plane. The expected exclusion limit is indicated by the blue dashed line, with its uncertainties given by the yellow band. The observed limit is indicated by the solid maroon line, while the dotted lines represent the limit when the signal expectation is deviated by its $\pm 1\sigma$ theoretical uncertainty [169]. The purple shaded area shows the observed limit from the 2011 $\sqrt{s} = 7 \text{ TeV}$ analysis, while theoretically excluded regions are shaded green and blue [262]. Figure (c) shows the expected limits for each of the ‘tight’ selections.

various production mechanisms are allowed.

The following chapter will discuss these results, while also discussion future prospects and implications of these results.

Discussion and implications

In the previous chapter we have set lower limits on allowed sparticle masses from a search for strongly interacting SUSY particles, using 5.8 fb^{-1} of data taken at $\sqrt{s} = 8 \text{ TeV}$. This chapter will discuss these results and put this result in perspective, by discussing implications and future prospects.

6.1 Discussion

Looking back at the results given in section 5.6, unfortunately no SUSY discovery can be claimed. While in the previous $\sqrt{s} = 7 \text{ TeV}$ analysis [169] the tight 4-jet signal region had one 2σ deviation, no such deviation is seen here, confirming the suspicion that the observed deviation was merely statistical.

However, large parts of the SUSY phase space have been excluded by the analysis, setting strong requirements on supersymmetry. Exclusion limits have been set on the constrained MSSM and three simplified models, which increase the excluded SUSY phase space in the case of light LSPs with respect to the $\sqrt{s} = 7 \text{ TeV}$ limits. For models containing an LSP with a mass close to the squark or gluino mass, the limits agree with these previous results. The exclusion limit in the CMSSM model has significantly improved by the increase in statistics and centre of mass energy, for which the analysis was re-optimised. Although in this scenario the masses of squarks and gluinos are limited to above 1.5 TeV and 1 TeV respectively (figure 5.29), this does not mean these mass limits are true for any general SUSY model, since the CMSSM has many assumptions – e.g. on the mass relations between the particles in the gaugino sector. For instance, the interpretation in the simplified MSSM models shows that if only pure squark-antisquark production would be possible at the LHC, the mass limit on squarks would be lowered to 780 GeV. For a high mass LSP no exclusion is possible, as seen in figure 5.25.

SUSY with a compressed spectrum remains difficult to reach. An important observation can be made when viewing the directly decaying squark-pair model: an optimised analysis for 5.8 fb^{-1} of $\sqrt{s} = 8 \text{ TeV}$ data could not improve the result of the 4.7 fb^{-1} $\sqrt{s} = 7 \text{ TeV}$ analysis. The same holds for the limit for gluino-pair production, where

the analysis presented in this thesis only improved the high mass region. This is mainly due to increased background predictions and increased systematic uncertainties in the looser selections. The source of the increase of systematic uncertainties is the uncertainty on W +jets events, coming from mismodelling of heavy-flavour quarks.

To be sensitive to the very compressed region, future analyses will need to lower the systematic uncertainties and introduce new variables or techniques. One such possible technique would be a multi-variate analysis specifically for the compressed region. At the same time, the systematic uncertainties need to be lowered to be able to benefit from any new technique. The main current sources of systematic uncertainties on the SM background are the heavy-flavour issue of the SHERPA W +jets sample, theoretical uncertainties and a low number of events in some CRs, while for signal events the ISR modelling is an important source of uncertainty. In a future analysis, the W +jets uncertainty could be reduced. With a larger dataset some of the statistical uncertainties in the CRs may decrease, leading to smaller uncertainties on the transfer factors. Improving the uncertainties from detector-related sources, such as the JES and $\cancel{E}_T$ uncertainties, will require a lot of effort from dedicated analyses, while they only contribute minimally. Since diboson production is becoming a significant background in the low jet multiplicity channels, the lack of a diboson control region is a big impact on the analysis. With the dataset used in this thesis, a dedicated control region would suffer from low statistics. However, without such a CR, uncertainties cannot be decreased via transfer factors, leading to large contributions on the total background uncertainty from the diboson prediction. Defining such a diboson control region should thus be a priority for a future analysis with a larger dataset.

6.1.1 Update with full 2012 dataset

After completion of the analysis presented in this thesis, the hadronic search has been updated and re-optimised for the full 20.3 fb^{-1} of data recorded in 2012. The updated analysis [263] builds on the analysis presented here. In the updated analysis, the signal region definition is similar to ours, with the most significant change in SRA ‘medium’. This signal region uses instead of a requirement on $\cancel{E}_T/m_{\text{eff}}^{\text{excl}}$ a cut on the $\cancel{E}_T$ significance, $\cancel{E}_T/\sqrt{H_T}$, with H_T the scalar sum of the p_T of all jets. The preliminary result is given in figure 6.1 for the gluino-pair and squark-pair simplified models. In both interpretations the limits have improved in the high mass region, excluding gluinos for $m_{\tilde{g}} < 1440 \text{ GeV}$ and squarks with $m_{\tilde{q}} < 840 \text{ GeV}$ at $m_{\tilde{\chi}_1^0} = 0 \text{ GeV}$. In the compressed region the limits are slightly better than those obtained with 5.8 fb^{-1} : the limit in the squark-pair model has increased around 50 GeV in LSP mass, while it increased $\sim 25 \text{ GeV}$ in LSP mass for the gluino-pair model. The main improvements come from the corrected SHERPA W +jets samples, which reduce the uncertainty on W +jets events, and from additional data statistics. Again, as with the analysis described in this thesis, without a diboson control region, the uncertainty on the diboson prediction is one of the leading uncertainties.

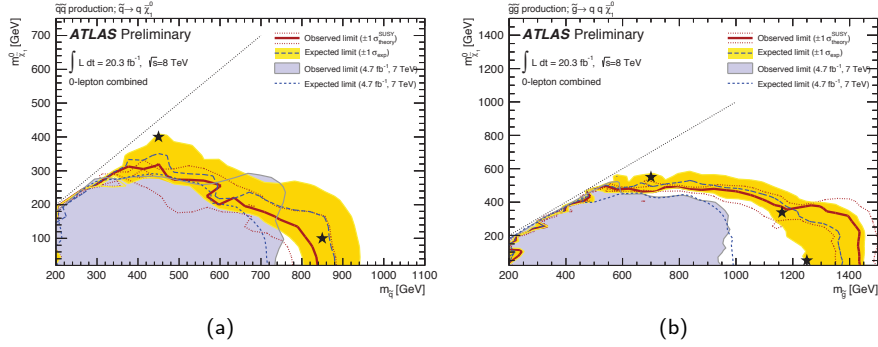


Figure 6.1: Exclusion limit for the direct production of (a) squark pairs and (b) gluino pairs, and their decay into quarks and lightest neutralinos, in 20.3 fb^{-1} of data taken at $\sqrt{s} = 8$ TeV, compared to the limits of the $\sqrt{s} = 7$ TeV analysis. The stars represent benchmark points used in the analysis paper. Figure taken from [263].

6.2 Implication of the results

What should we take away from the results of the presented analysis? Looking at figures 5.25-5.29, we observed that squarks and gluinos are excluded up to 1500 GeV for equal mass squarks and gluinos in the CMSSM, while in the simplified models we find $m_{\tilde{q}} < 780$ GeV for squark-pair production, $m_{\tilde{g}} < 1440$ GeV for the $\tilde{q}\tilde{q}$ model and $m_{\tilde{g}} < 1175$ GeV for gluino-pair production. As already noted before, the exclusion limits on SUSY production degrade quickly if the LSP is close in mass to the squark or gluino due to there being less missing transverse momentum and softer jets in decays when SUSY particles are produced in these models.

This is an important observation. Even though gluinos are excluded up to 1.2 TeV for massless neutralinos, we do not have any limit on the gluino mass for $m_{\tilde{\chi}_1^0} = 500$ GeV in the case of gluino-pair production. The compressed region with the LSP mass near either the gluino or (one of) the squark masses remains a possible hiding place for SUSY with gluino/squark masses below 1 TeV.

Another important statement should be made about model-dependency. The difference between the CMSSM limit and simplified models alone shows that these are not universal lower limits, but depend on the specific model. Since SUSY is a landscape full of possible models, with either R-parity conservation or violation, with many different possibilities for SUSY breaking, and where the minimal SUSY model has 105 parameters, one should take care when interpreting results. To be sure about the feasibility of a particular SUSY model, one should simulate events and emulate the analysis cuts. For this goal the introduction of simplified models is an important aid, since theorists can validate their implementation of our analysis by comparing to our efficiency for pure production of for instance gluino-pairs.

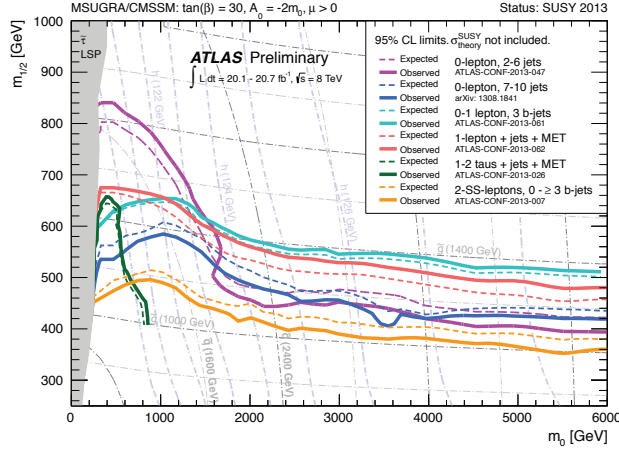


Figure 6.2: Comparison between different ATLAS SUSY searches of observed (solid lines) and expected (dashed) exclusion limits. All searches are based on $\sim 20 \text{ fb}^{-1}$ of $\sqrt{s} = 8 \text{ TeV}$ data.

6.2.1 Results from other search channels

To be sensitive over the full range of (R-parity conserving) SUSY models, many different searches are defined, each targeting a different signature in the detector. Besides the hadronic search, searches for the production of strongly interacting SUSY particles include, amongst others: a jets plus $\cancel{E}_T$ search selecting one or more leptons [264]; an analysis targeting events with at least 3 b -jets and $\cancel{E}_T$ [265]; and an analysis selecting events with at least 7 jets and $\cancel{E}_T$, without leptons [266]. To find out what the current status of SUSY is, given the results of these searches, we can proceed in a number of ways, for instance by comparing them in a specific model, or by studying the general implications of the searches on the MSSM. The former is discussed in this section, while the next section discusses shortly an analysis on a general framework.

Although constrained models do not give a perfect overview of the status of SUSY as a whole, the CMSSM can be used to compare various ATLAS searches. The production and subsequent decay of SUSY particles leads to different signatures depending on the choice of parameters in the CMSSM. For low m_0 , the production cross section is dominated by squark-pair production, decaying either directly or via charginos into neutralinos. The hadronic search is thus expected to do best in this region, together with a search for high $\cancel{E}_T$ with leptons. However, for high m_0 and low $m_{1/2}$, gluino-pair production dominates. Since in this region gluinos are lighter than 1st/2nd generation squarks, but heavier than stops, the gluinos will often decay into a stop-top pair, with the stop decaying further into either a top or a bottom, plus a neutralino. In this region searches targeting b -jets or large jet multiplicities will perform well.

Figure 6.2 shows a comparison between (amongst others) the analyses discussed

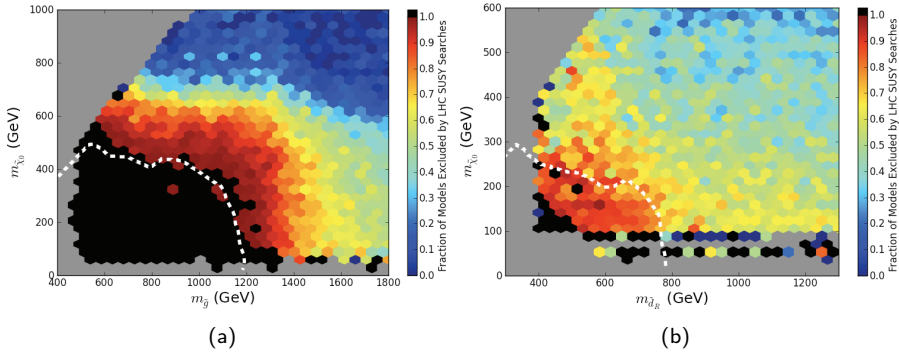


Figure 6.3: Fraction of pMSSM models excluded by LHC searches, projected on a (a) $(m_{\tilde{g}}, m_{\tilde{\chi}_1^0})$ and (b) $(m_{\tilde{a}_R}, m_{\tilde{\chi}_1^0})$ mass plane. The dashed white lines represent the limits set by the analysis described in this thesis. Figure taken from [267].

above of the observed (solid lines) and expected (dashed) exclusion limits in the CMSSM. All the shown analyses are using $\sim 20 \text{ fb}^{-1}$. The update of the hadronic search discussed in section 6.1.1 is the best performing analysis for low m_0 , while the analysis targeting at least 3 b -jets sets the most stringent limit for high m_0 .

6.2.2 Implications of the analysis on a phenomenological MSSM

Results set in constrained models such as the CMSSM are very model-dependent. To be less model-dependent, we have previously introduced the simplified models. A third approach followed by various theorists is to interpret the LHC results in a more general manner. One such effort, described in refs. [267, 268], has an interesting approach: they reduce the full MSSM to the most general R-parity conserving 19-dimensional parameter space by making a number of assumptions and requiring the MSSM to obey various experimentally-motivated principles. In the resulting *phenomenological* MSSM, or pMSSM, hundreds of millions of points are randomly picked for which MC simulated events are generated. The generated points are required to have squark and gluino masses of less than 4 TeV and have $\tilde{\chi}_1^0$ as the LSP.

The generated events are passed through an emulation of a large set of ATLAS (and some CMS) SUSY analyses. To study which pMSSM points are excluded and which survive the current LHC limits the results are compared to the exclusion limits set by the experiments. The implemented analyses are most ATLAS SUSY 7 and 8 TeV analyses, supplemented by some CMS analyses on $\sqrt{s} = 7 \text{ TeV}$ and the measurement of $B_s \rightarrow \mu^+ \mu^-$. From this comparison some general conclusions can be derived.

From the 225000 points surviving first constraints from theoretical, dark matter and flavour physics sources, 37% are excluded from the combination of LHC searches, while the hadronic search presented in the previous chapter excludes 27% on its own. Figure 6.3 shows the fraction of pMSSM points excluded by the combination of LHC

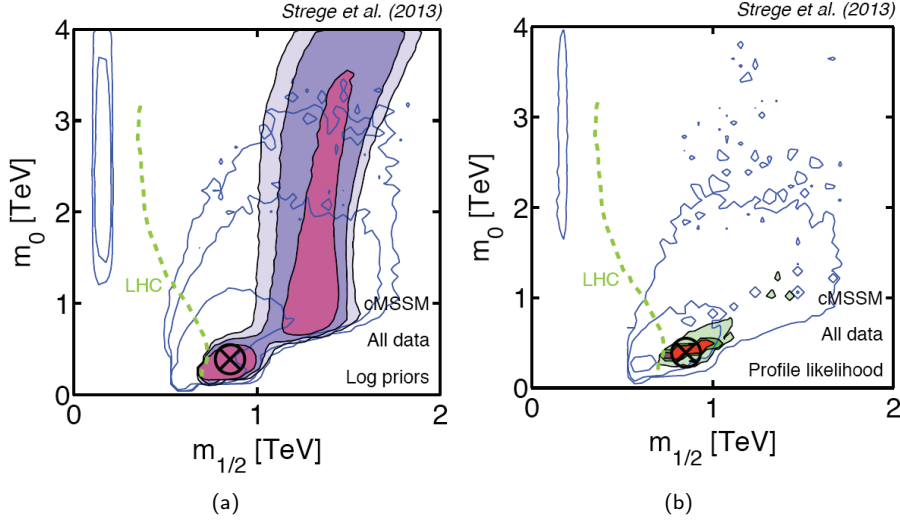


Figure 6.4: Constraints from the on the CMSSM $(m_{1/2}, m_0)$ plane, using ATLAS SUSY results on 5.8 fb^{-1} of $\sqrt{s} = 8 \text{ TeV}$ data, the Higgs mass measurement, the $(g - 2)_\mu$ constraint and the XENON100 2012 results. The constraints are set using Bayesian methods with logarithmic priors (a) and a frequentist one-dimensional profile-likelihood (b). The filled contours represent the 68%, 95% and 99% confidence regions, while the cross marks the point with the best fit overall. The open blue contours represent the previous results using only 1 fb^{-1} ATLAS results. The green line represents the ATLAS exclusion limit. Note that m_0 and $m_{1/2}$ are interchanged with respect to other CMSSM representations. Figure taken from [269].

searches, projected on a gluino-LSP mass plane (left) and squark-LSP mass plane (right), where in the latter case the right-handed down-type squark is chosen. The dashed white line represents the exclusion limit from the hadronic search.

From figure 6.3 (a) one can conclude that the gluino limit set by our hadronic search on a simplified model with pure gluino-pair production is actually a good indication of the limit in the whole of the MSSM. However, note that the figure shows the fraction of excluded models by the combination of searches, not just the hadronic search. Further studies show that the analysis selecting events with at least 3 b -jets contributes to this fraction for gluinos decaying to stop-top pairs. As was observed before, less models with a high LSP mass are excluded.

For light squarks the situation is different. In our simplified models, the assumption is made that the left- and right-handed squarks of the first two generations are mass degenerate. However, this needs not be the case in the MSSM. Figure 6.3 (b) shows the fraction of excluded models as a function of the $m_{\tilde{d}_R}$ and $m_{\tilde{\chi}_1^0}$ mass. The limit

set by us in the squark-pair simplified model does not correctly represent the excluded fraction for these right-handed squarks. In general, the production cross section is smaller for $\tilde{d}_R$, and due to smaller PDFs for d -quarks than for u -quarks. Moreover, in the pMSSM the mass matrices for fermions are diagonal leading to degenerate $m_{\tilde{u}_L}$ and $m_{\tilde{d}_L}$, while $m_{\tilde{u}_R}$ and $\tilde{d}_R$ are not correlated. Therefore, models with a light $\tilde{d}_R$ may have other heavier squarks, leading to a smaller squark-pair cross section. As a result, models with light right-handed squarks may still survive all ATLAS SUSY searches.

6.2.3 General SUSY fits

Finally, the effect of LHC searches on SUSY can also be studied from global fits to data from various measurements in specific models, in particular the CMSSM. These global fits search for the region of the CMSSM which satisfies the constraints best. Several collaborations, such as FITTINO [270], MASTERCODE [271], BAYESFITS [272] and Strege et al. [269], have implemented ATLAS and CMS searches for SUSY using jets plus $\cancel{E}_T$ conducted with the full $\sqrt{s} = 7$ TeV dataset, as well as dark matter detection experiments, $(g - 2)_\mu$ results and flavour physics results. No other SUSY searches have been included, as the hadronic search is the most sensitive [271]. The last mentioned reference is the only one currently using 2012 LHC datasets for SUSY limits, the $B_s \rightarrow \mu^+ \mu^-$ measurement and the Higgs mass measurement.

Figure 6.4 shows the contour lines of the fit by Strege et al. for fits using a Bayesian technique with priors uniform in the logarithm of the masses (left) and using a frequentist profile likelihood technique. Both techniques show that the $\sqrt{s} = 8$ TeV results of the hadronic SUSY search reduce the 68% CL contours to a small regions at $800 < m_{1/2} < 1000$ and $300 < m_0 < 400$, where for the Bayesian technique there is also a large area at high m_0 available. The $(g - 2)_\mu$ constraint requires low values for m_0 and $m_{1/2}$, while conversely the ATLAS searches require higher masses, leading to friction between the constraints. Ignoring the $(g - 2)_\mu$ constraint would leave more room for SUSY at high m_0 and $m_{1/2}$ in the CMSSM. The point which satisfies the constraints best lies at $m_0 = 389$ GeV, $m_{1/2} = 853$ GeV, $A_0 = -2664$ GeV and $\tan \beta = 14.5$. This point lies very close to the 20 fb^{-1} exclusion line, and thus the low m_0 best fit region is within reach of the LHC in the coming years. However, as the goodness of fit is deteriorating from difficulties to satisfy all the experimental constraints, the CMSSM seems to be disfavoured, and more general SUSY scenarios should really be studied, as for instance with the pMSSM described before.

6.3 Outlook for $\sqrt{s} = 13$ TeV

When the LHC restarts in 2015 after a two-year-long technical shutdown, the centre-of-mass energy will have increased to 13 TeV, while instantaneous luminosity will rise to $\sim 10^{34} \text{ cm}^{-2}\text{s}^{-1}$. The projected integrated luminosity is between 75 and 100 fb^{-1} until the next long shutdown in 2018. At these energies, the cross sections for pair production will rise considerably in the simplified models, by a factor of 200 to near 1 fb

for 2 TeV gluinos, and a factor of 300 for 2 TeV squarks. With the four-fold increase of statistics, this increase in cross section will allow for a sensitivity for gluinos and squarks with a mass up to 2 TeV.

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Summary

Ever since I started studying physics, I have heard the sentence ‘as a high school student I always thought physics was fascinating, but I just don’t understand any of it’ more often than I can count. During my years working as a PhD student in particle physics the interest but also the incomprehension only increased among my friends and family. Therefore in the first pages of this summary I will attempt to shortly explain the concepts behind particle physics.

My contributions to particle physics are in the field of supersymmetry. This theory is introduced to solve several issues with the current theory of particle physics, called the *Standard Model*. I have searched for signs of supersymmetry in particle collisions in the ATLAS detector at CERN. The second part of this summary will elaborate on the experimental analysis performed and documented in this thesis.

Particle physics and the Standard Model

The subject of particle physics attempts to explain how nature works on the smallest currently testable length scales. This basically comes down to answering the questions ‘what are we (and everything around us) made of?’ and ‘how do the smallest building blocks interact with each other’?

All matter on earth is made of molecules composed of atoms. For example, water molecules consist of two hydrogen atoms and one oxygen atom. Since the beginning of the 20th century it is known that atoms are not ‘indivisible’, as their Greek name suggests, but consist of a nucleus made from protons and neutrons with electrons surrounding it. In the 1960s it was discovered that even protons and neutrons are composite particles: they are composed of a combination of *quarks*. These quarks, together with several other particles such as the electron, are finally currently considered to be the most fundamental particles. The theoretical model describing all fundamental particles is known as the Standard Model of particle physics. It also describes three of the four fundamental forces between particles: electromagnetism, the strong and the weak nuclear force. Electromagnetism is well known from daily life, for instance from electromagnetic radiation (e.g. light) and is the driving force behind most chemistry processes. The other two are less well known, since they only act such short length scales that we do not experience them consciously. The strong nuclear force binds protons and neutrons together in the atomic nucleus. The weak nuclear force is the only force which can change the flavour of quarks, and is thus responsible for radioactive decay of subatomic particles. Gravity is not incorporated in the theory, since the large scales on which it acts cannot be united with the underlying quantum

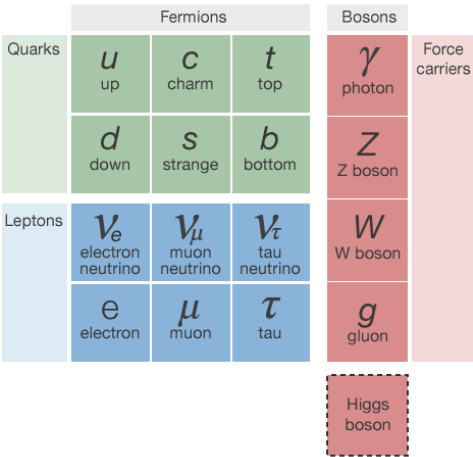


Figure S.1: Graphical representation of the particles in the Standard Model. Figure taken from [273].

mechanics of the Standard Model.

To make sense of the big jumble of particles in the Standard Model, they are categorised according to their properties, as illustrated in figure S.1. Particles with a half-integer *spin*, a quantum-mechanical form of angular momentum, are called *fermions*. Among the fermions are six flavours of *quarks*, while the electron is an example of one of the six *leptons*. Fermions are generally considered to be matter particles. Everything we know is built up of only the lightest family of fermions (up and down quarks, electrons and electron neutrinos), since heavier fermions are unstable: when produced, they decay after a very short time into the lighter stable particles.

Interactions between the fermions are mediated by particles with an integer spin, called *bosons*. For example, the electromagnetic force is mediated by photons, particles which we perceive as light. The decay of heavy particles often involves gluons, *W* and *Z* bosons. The former mediate the strong force, while the latter two mediate the weak force. The Higgs boson, recently found at CERN, is a special particle, since it does not really mediate a force, but is a ‘collateral’ particle coming from the Brout-Englert-Higgs mechanism through which particles obtain mass.

Shortcomings of the Standard Model

Throughout the 20th century, the Standard Model described above has been tested to incredible precision. However, for all its success, there are several shortcomings to the model causing the physics community to doubt that it is the final theory of fundamental particles.

One of these shortcomings has to do with cosmological observations of additional

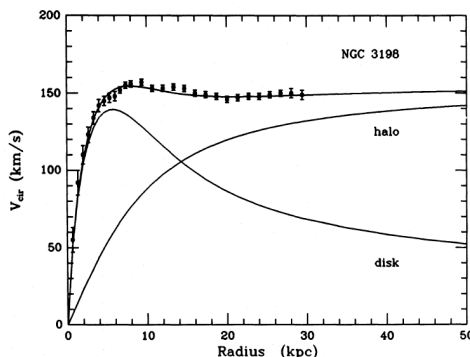


Figure S.2: *Measured stellar orbital velocity of stars in spiral galaxy NGC 3198, given as a function of distance to the galactic centre is shown by the black blocks. The lower curve shows the expected rotational curve assuming no dark matter, with all visible matter in a disk, while the middle curve shows the dark matter contribution to fit the data [48].*

matter of unknown origin in the universe. The first sign of such matter was observed in the 1930s from measurements of the orbital velocity of stars in spiral galaxies. Considering the assumed particle content in the galaxy and Newtonian laws of gravity, it is expected that the orbital velocity for stars decreases as the square root of the distance to the galactic centre, as illustrated by the lower line in figure S.2. However, observations show that the orbital velocity remains nearly constant. The most popular explanation is that there exists a halo of unknown matter around galaxies. This *dark matter* does not (or only very weakly) interact with the 'normal' matter inside the galaxy in any other way than via gravity. Without interactions the matter does not heat up and is not visible to us. Because the effect is not consistent with any known particles, dark matter should consist of some unknown particle(s). Since this observation the evidence for the dark matter hypothesis has increased enormously, for instance via measurements of the cosmic microwave background radiation. It is estimated that only 15% of all matter in the universe is in the form of ordinary Standard Model matter, while the remaining 85% is of the form of dark matter.

Besides dark matter, there are several theoretically reasons to believe something is missing in the Standard Model, one of which is called the *hierarchy problem*. It states that there is a conflict between energy scales in particle physics: on the one hand the mass of the Higgs boson is required to be low for the Brout-Englert-Higgs mechanism to do its job, while on the other hand quantum-mechanical effects in its theoretical calculation allow the Higgs mass to be enormous. To counter the effect, bringing the calculated Higgs mass back to the preferred lower mass scale, involves fine-tuning of the calculation by a factor of 10^{17} , which does not feel right - it would mean the reason this universe works as it does is just because two gigantic numbers nearly cancel each other by chance.

Supersymmetry

Since the Standard Model has been so successful in its predictions, one way to deal with the above issues is to use the Standard Model as a foundation for a more complete model, instead of moving to a new theory entirely. This new model should describe the Standard Model as well as ‘new physics’, allowing for solutions to the shortcomings listed above. One such theory introducing new physics is called *supersymmetry* (SUSY). As the name suggests, it is based on a fundamental symmetry in nature. In particle physics a symmetry should be understood as a transformation under which the physical laws remain unchanged.

The Standard Model incorporates many such symmetries, however it is still possible add one: we can require that bosons and fermions can be interchanged without changing the physics of particle interactions. Including this symmetry requirement in the Standard Model leads to supersymmetry: the hypothesis that for each boson there is a fermion with the exact same properties except for its spin, and likewise for each fermion there is such a boson. Since these (supersymmetric) partner particles have not yet been observed in nature they should be heavier than their Standard Model partners, which is possible by breaking the symmetry at a high energy. One of these additional particles is a perfect dark matter candidate: the neutralino $\tilde{\chi}_1^0$. Supersymmetry also allows for a solution to the hierarchy problem. However, since there are many ways of breaking the symmetry, and since the masses of the SUSY particles are unknown, we do not have one single SUSY theory, but are left with a vast landscape of possible SUSY scenarios.

Goal of this thesis

With the introduction of supersymmetry, we have arrived at the goal of this thesis. Since we know the Standard Model is not the final theory of particle physics, we want to detect hints of physics beyond the Standard Model. To this end, together with a large group of physicists at CERN, I have searched in proton-proton collisions for deviations from the Standard Model expectations consistent with supersymmetry. If large deviations are observed it can lead to the discovery of SUSY, giving a whole new view on physics in the early universe and raising many new tantalising questions. On the other hand, even though we cannot disprove supersymmetry completely, observing no significant deviations is very useful as well. In this case, we can set exclusion limits. These give the upper limit on masses of SUSY particles for which SUSY is excluded, while above these limits we cannot rule it out. It will give us a lot of information on the remaining possibilities for SUSY to exist, and will guide future physicists in their search for new physics.

Detecting particles

To search for supersymmetric particles, one needs to create an environment to both create them and (indirectly) detect them. Since most SUSY particles are unstable, they do not exist in ordinary matter. To study or discover them, they first need to

be produced by colliding protons at very high energy at the Large Hadron Collider (LHC) at CERN in Geneva. Einstein's famous $E = mc^2$ tells us that energy can be transformed into mass and vice versa. Therefore, when two particles with high energies collide, one or more heavier particles can be produced. The LHC is currently the world's largest and most powerful particle accelerator, colliding protons at a centre-of-mass energy of 8 TeV, where 1 TeV is equivalent to the energy of a flying mosquito, which consists of the order of 10^{23} atoms.

After the production of supersymmetric or other heavy particles (such as top-quarks), these will decay immediately into lighter quarks and leptons. Although these particles themselves are too small to be perceived directly, they can be indirectly detected using the ATLAS particle detector, situated around a collision point of the LHC. Particles travelling through the detector ionise certain material and are slowed down and stopped in calorimeters, via which the particles can be reconstructed and their momentum measured.

Searching for supersymmetry

Using data recorded from proton-proton collisions, or *events*, by the various detector elements of ATLAS I searched for signs of supersymmetric particles in proton-proton collisions. To this end I used a 'counting experiment'. Each event is categorised as being either Standard Model-like or SUSY-like. The observed number of SUSY-like events is compared to the number expected from collisions involving only Standard Model particles.

However, each event comes with a calculable probability of producing certain particles, which depends on the mass of the involved particles. Since SUSY particles are expected to be heavier than their Standard Model counterparts, in each event there is a smaller probability of producing SUSY particles than 'normal' Standard Model particles. In the 20 million events occurring per second many more Standard Model particles are produced than (hypothetical) SUSY particles. The selection of SUSY-like events therefore needs to be carefully optimised.

Selection of the hadronic channels

To differentiate between SUSY and Standard Model events it is essential to accurately know the signature which either type of event leaves in the detector. In large portions of the possible SUSY scenarios the majority of events lead to 2 strongly interacting SUSY particles being produced that decay into high-momentum quarks and two neutralinos. Solitary quarks are not allowed in the Standard Model, causing the outgoing quarks to hadronise into a shower or *jet* of particles. Neutralinos are stable particles which do not interact with the detector and are therefore invisible to us. However, due to momentum conservation, we do observe a net loss of momentum in the transverse plane: since the two incoming protons have no transverse momentum, the outgoing particles should have zero net transverse momentum as well. If any particle, such as a neutralino, escapes detection, it will show up as a momentum imbalance, measured as missing transverse energy $\cancel{E}_T$.

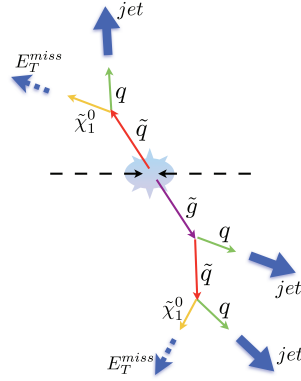


Figure S.3: Illustration of the production and immediate decay of a squark and a gluino into quarks and the lightest supersymmetric particles, neutralinos ($\tilde{\chi}_1^0$), leading to a signature of jets and missing transverse energy.

Such SUSY scenarios thus have a *hadronic* signature of two or more jets and high missing transverse momentum, while leptons are less likely in the decays. Such a decay is illustrated in figure S.3. To reduce the Standard Model backgrounds to our search, we select events with high $\cancel{E}_T$ and high momentum jets. Since the decay of most Standard Model particles involves leptons, requiring no leptons in our search reduces the backgrounds even further. Finally, the selections are optimised for different types of SUSY using information on $\cancel{E}_T$, the number of jets and their momentum. All in all 12 such signal selections are defined, which all have a final requirement on the *effective mass* m_{eff} of the event. This is defined as the sum of $\cancel{E}_T$ and the transverse momenta of the jets in the event. The effective mass is correlated to the mass of the SUSY particles, and will thus be higher on average in SUSY events than in Standard Model events.

Background estimation

Once the signal selection criteria have been established, we need to know how many data events are selected if SUSY does not exist. This Standard Model background is estimated using a combination of simulations of the data and extrapolations from Standard Model rich selections to our SUSY rich selection. The simulations are performed by Monte Carlo event generators. These simulate events from theoretical principles, including the production of particles in a collision, as well as their decay and the interaction of the decay particles with the detector. It is essential to know and reduce the uncertainties on the background estimates. The uncertainties have both a theoretical origin, due to the used simulations, and an experimental origin, due to detector calibrations.

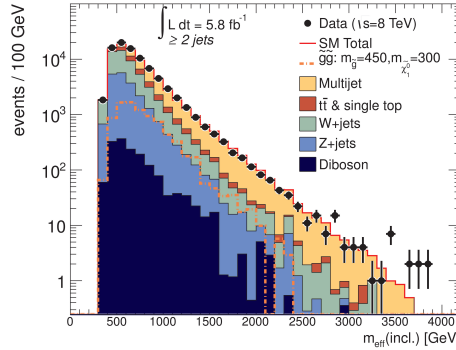


Figure S.4: Observed m_{eff} distribution with the final selection requirement indicated by the red arrow. Data is given by black dots, while the Standard Model backgrounds are illustrated by the stacked coloured histograms. The dashed histogram represent a SUSY model.

Supersymmetric particle production

Besides these background estimates, the number of SUSY signal events in the selection also needs to be estimated. To this end, events are simulated with supersymmetric particle production. Knowledge of the number of hypothetically produced SUSY particles is essential. The production rate of any fundamental particle can be calculated using its *cross section*. The calculation of the cross section depends on several theoretical and experimental inputs, of which *parton distribution functions* (PDFs) are the most important. They are used to estimate the momentum-dependent probability that for instance a quark in one incoming proton collides with a gluon in the other proton. However, since PDFs are difficult to measure, they carry large uncertainties which propagate into an uncertainty on the value of the calculated cross section. Using methods I have developed it is determined that this propagation leads to uncertainties on the production rate of *squark* and *gluino* pairs, the SUSY partners of quarks and gluons of between 10% for low mass particles and 70% for very massive particles. Since these theoretical uncertainties only affect the signal, their effect is shown separately from the background uncertainties in the following results

Results

Only once all the above-mentioned groundwork has been done, the data taken in the first half of 2012 is tested against our selections. In figure S.4 the number of selected events is plotted against their observed m_{eff} for a signal selection where at least two jets are selected. The black dots represent the data points, while the stacked coloured histograms represent the various estimated backgrounds. The dashed line represents a m_{eff} distribution for a SUSY model. If SUSY exists with squarks and gluinos with such masses, large discrepancies would be seen between data and expectations at high values of m_{eff} , to the right of the red arrow. However, no significant deviation is seen

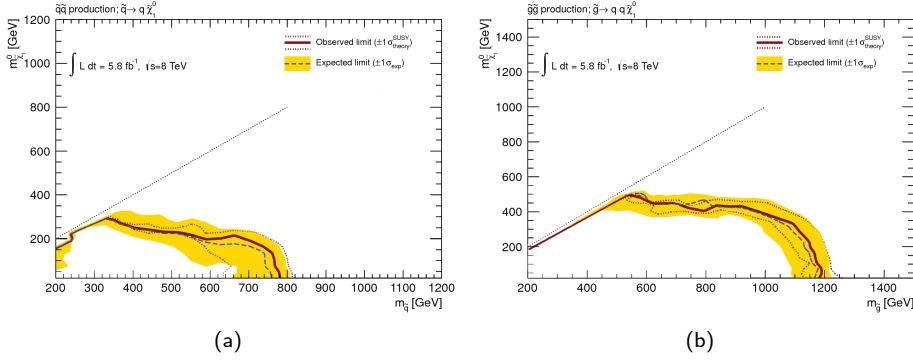


Figure S.5: Exclusion limits for producing two squarks (a) or two gluinos (b). SUSY models are excluded if the squark or gluino mass and neutralino mass fall inside the excluded area. The red line indicates the observed limit, while the blue line indicates the expected limit if data would follow the background estimates exactly. The yellow band gives the uncertainty on the expected limit.

in this or any other selection. With no large observed discrepancies, the number of possible supersymmetric signal events hiding in the data is calculated using a statistical framework, based on the theoretical and experimental uncertainties on the background estimates. This number can tell us if certain supersymmetric scenarios can be ruled out. As mentioned before, more massive particles are less frequently produced, thus for very high masses we will not be able to draw any conclusions.

Since there are many possible SUSY scenarios, we need to choose in which we interpret our results. To be fairly model independent, the excluded masses are given for two hypothetical *simplified* SUSY scenarios where a restriction is set on which particles may be produced. On the left hand side in figure S.5 a the limit is shown in a model shown in which only squarks are produced, while the right figure shows a model with only gluino production. Although these simplified models cannot be the real manifestation of SUSY, if combined limits set in these can be used to configure a limit in any type of SUSY model. The limits are given as function of either the squark or gluino mass (x -axis) and neutralino mass (y -axis). The red line indicates the observed limit: SUSY models with squark or gluino masses below this limit are ruled out. The dotted red lines indicate this same limit when the cross sections for SUSY production are taken to be consistent with their (up- or downward) uncertainty discussed previously. In SUSY with massless neutralinos, squarks are limited to be heavier than $\sim 780 \text{ GeV}^1$, while gluinos should be heavier than $\sim 1175 \text{ GeV}$. However, the limits depend greatly on the neutralino mass, into which squarks and gluinos decay. For very heavy neutralinos, with masses close to the squarks and gluinos, SUSY cannot be excluded.

¹if all partners of the five lightest quarks are of equal mass

Conclusions and outlook

Although unfortunately no indications of supersymmetry have yet been found the analysis in this thesis only scratches the surface of the SUSY landscape, leaving many possibilities for SUSY to still exist. One such possibility is for it having neutralino masses close to that of squarks and gluinos. However, it is being hunted down and slowly cornered² due to the relentless and ongoing efforts of many physicists around the world. These efforts will only increase in the coming years, when the LHC will reach higher collision energies and record up to one hundred times more data. Even if no evidence for SUSY is found, I am hopeful that such searches at the LHC will be expose anomalies in the Standard Model.

²At least the most interesting SUSY models with not too heavy particles.

Samenvatting

Vanaf het moment dat ik natuurkunde ben gaan studeren heb ik de zin 'ik vind natuurkunde altijd heel interessant, alleen begrijp ik er nooit iets van' vaker gehoord dan mij lief is. In de afgelopen vier jaar waarin ik promoveerde in de deeltjesfysica is deze interesse vanuit mijn omgeving alleen maar gegroeid, maar helaas ook het bijbehorende gevoel van onwetendheid. Ik zal daarom de eerste pagina's van deze samenvatting gebruiken om kort de concepten achter deeltjesfysica uit te leggen.

Mijn onderzoek vond plaats binnen de deeltjesfysica, met als onderwerp supersymmetrie. Deze theorie is geïntroduceerd om verschillende problemen op te lossen die het huidige *Standaard Model van deeltjesfysica* ondervindt. Ik heb gezocht naar aanwijzingen van supersymmetrie in deeltjesbotsingen in de ATLAS detector op CERN. Het tweede gedeelte van deze samenvatting zal dieper ingaan op de uitgevoerde experimentele analyse die in dit proefschrift is gedocumenteerd.

Deeltjesfysica en het Standaard Model

Deeltjesfysica heeft ten doel uit te leggen hoe de natuur werkt op de allerkleinste (op dit moment toetsbare) afstanden. Dit komt uiteindelijk neer op het beantwoorden van vragen als 'waar zijn wij (en alles om ons heen) van gemaakt?' en 'hoe hebben deze aller kleinste bouwblokken interactie met elkaar?'.

Alle aardse materie bestaat uit moleculen, samengesteld uit atomen. Watermoleculen bestaan bijvoorbeeld uit twee waterstof atomen en een zuurstof atoom. Sinds het begin van de 20^e eeuw weten we dat atomen helemaal niet 'ondeelbaar' zijn, zoals hun Griekse naam suggereert, maar opgebouwd zijn uit een atoomkern bestaand uit protonen en neutronen waar een of meer elektronen omheen cirkelen. In de jaren zestig van de 20^e eeuw is ontdekt dat protonen en neutronen zelfs deelbaar zijn in kleinere deeltjes: ze zijn opgebouwd uit combinaties van *quarks*. Het zijn deze quarks, samen met elektronen en soortgelijke deeltjes, die tegenwoordig worden beschouwd als de meest fundamentele deeltjes. Het theoretische model dat deze deeltjes beschrijft is het Standaard Model van deeltjesfysica. Het beschrijft ook drie van de vier fundamentele krachten tussen deze deeltjes: elektromagnetisme, de sterke en de zwakke kernkracht. Elektromagnetisme is zeer bekend uit het dagelijks leven, bijvoorbeeld van elektromagnetische straling (licht), en is de drijvende kracht achter de meeste chemische processen. De andere twee zijn minder voor de hand liggend, aangezien ze alleen op zulke kleine afstanden werken dat wij ze niet bewust ondervinden. De sterke kernkracht bindt protonen en neutronen samen in een atoomkern. De zwakke kernkracht is de enige kracht die de 'smaak' van quarks, of de soort, kan veranderen. Hiermee is

het verantwoordelijk voor radioactief verval van subatomaire deeltjes. Zwaartekracht wordt niet door het Standaard Model beschreven aangezien de grote afstand waarop het werkt niet kan worden verenigd met de onderliggende kwantummechanica.

Om niet de weg kwijt te raken in de grote hoeveelheid deeltjes van het Standaard Model zijn ze ingedeeld naar hun eigenschappen, zoals weergegeven in figuur S.1. Deeltjes met een halfvallige *spin*, het kwantummechanisch equivalent van impulsmoment, noemen we *fermionen*. Onder deze fermionen zijn zes smaken quarks, terwijl het elektron een van de zes *leptonen* is. Fermionen worden beschouwd als materie-deeltjes. Alles om ons heen is opgebouwd uit fermionen, maar wel alleen van de lichtste familie: de *up* en *down* quarks, elektronen en elektron neutrino's. Zwaardere fermionen zijn instabiel: nadat ze gemaakt zijn vervallen ze gelijk weer in de lichtere stabiele deeltjes.

Interacties tussen fermionen worden overgebracht door *bosonen*, deeltjes met een geheelvallige spin. Zo wordt de elektromagnetische kracht overgebracht door fotonen, ook wel bekend als 'lichtdeeltjes'. Het verval van zware fermionen gaat via gluondeeltjes, *W* en *Z* bosonen. Gluonen brengen de sterke kernkracht over, terwijl *W* en *Z* bosonen de zwakke kernkracht overbrengen. Het nu zeer bekende Higgs boson, recentelijk ontdekt op CERN, is een speciaal deeltje, aangezien het niet zozeer een kracht overbrengt maar een bijkomend deeltje is van het Brout-Englert-Higgs mechanisme, dat deeltjes hun massa geeft.

Gebreken van het Standaard Model

Gedurende de 20^e eeuw is het Standaard Model uitgebreid getest tot ongelooflijke precisie. Echter ondanks dit succes zijn er enkele gebreken aan het licht gekomen waardoor de natuurkundige gemeenschap er vanuit gaat dat er nog een onderliggende theorie zal zijn.

Een van deze gebreken heeft te maken met de kosmologische observatie van materie in het heelal die niet in het Standaard Model is thuis te brengen. De eerste aanwijzing voor zulke materie werd gevonden in de jaren dertig van de vorige eeuw uit en komt uit metingen van de omloop snelheden van sterren in spiraalvormige sterrenstelsels. Met de huidige aannames van het deeltjesgehalte van sterrenstelsels en Newtons wetten van de zwaartekracht is de verwachting dat de omloopsnelheid van sterren zal afnemen met de wortel van de afstand tot het centrum. Dit wordt geïllustreerd met de onderste lijn in figuur S.2. Waarnemingen laten echter zien dat de omloopsnelheid bijna constant is. De populairste verklaring hiervoor is dat er zich een halo van onbekende materie om sterrenstelsels heen bevindt. Deze *donkere materie* heeft (zo goed als) geen interactie met 'normale' materie in het sterrenstelsel behalve via zwaartekracht. Zonder interacties wordt de materie niet warm en is dus niet zichtbaar voor ons. Aangezien het effect niet consistent is met voor ons bekende deeltjes zal donkere materie moeten bestaan uit nieuwe deeltjes. Na deze eerste aanwijzing is het bewijs voor de donkere materie hypothese alleen maar gegroeid, bijvoorbeeld via metingen van de kosmische achtergrondstraling. De huidige metingen geven aan dat slechts 15% van de materie in het heelal uit gewone Standaard Model deeltjes bestaat, terwijl de resterende 85% donkere materie is.

Behalve donkere materie zijn er enkele theoretische redenen om te denken dat er

iets mist in het Standaard Model. Een hiervan heet het *hiërarchie probleem*. Het stelt dat er een conflict is tussen energieschalen in deeltjesfysica: aan de ene kant moet de massa van het Higgs boson laag zijn om het Brout-Englert-Higgs mechanisme zijn werk te laten doen, terwijl aan de andere kant kwantummechanische effecten in zijn theoretische berekening de Higgs massa vrij laten om enorm groot te zijn. Om dit effect tegen te werken en zo de berekende Higgs massa terug te brengen naar de vereiste lage massa is er een zeer precieze afstemming nodig van constanten in de berekening, die tot op een factor 10^{17} goed gekozen moeten worden. Dit voelt gek: het betekent dat de reden dat het universum werkt zoals het doet is omdat toevallig twee gigantische getallen elkaar bijna (maar net niet) wegstrepen!

Supersymmetrie

Aangezien het Standaard Model zo succesvol is geweest met zijn voorspellingen, is het voor de hand liggend om de bovenstaande gebreken op te lossen met het Standaard Model als fundament, in plaats van deze in de prullenbak te werpen. Het nieuwe model moet naast het Standaard Model ook 'nieuwe fysica' beschrijven, waarmee de hierboven beschreven gebreken op te lossen zijn. Een populaire theorie die nieuwe fysica introduceert is *supersymmetrie*, of SUSY. Zoals de naam doet vermoeden is het gebaseerd op een fundamentele natuurkundige symmetrie. In deeltjesfysica wordt onder een symmetrie een transformatie verstaan waaronder de fysische wetten onveranderd blijven.

Het huidige model omvat veel van zulke symmetrieën. Er is er echter een die we nog zouden kunnen eisen: het verwisselen van bosonen en fermionen bij het beschrijven van een botsing zonder dat het een andere uitkomst geeft. Het Standaard Model met deze symmetrie daarin opgenomen noemen we supersymmetrie. De hypothese is dat voor elk boson er een fermion moet bestaan met exact dezelfde eigenschappen, met uitzondering van de spin, en omgekeerd voor fermionen. Omdat deze deeltjes nog nooit zijn waargenomen, zullen ze zwaarder moeten zijn dan hun Standaard Model partners, wat mogelijk is door de symmetrie te 'breken'. Een deeltje uit deze toegevoegde deeltjes is een perfecte kandidaat voor donkere materie: het neutralino $\tilde{\chi}_1^0$. Tegelijkertijd biedt supersymmetrie ook een oplossing voor het hiërarchie probleem. Omdat er echter verschillende manieren zijn om de symmetrie te breken en de massa's van de SUSY deeltjes onbekend zijn is er niet een enkele SUSY theorie, maar is er een grote verscheidenheid aan mogelijke SUSY scenario's.

Doel van dit proefschrift

Met de introductie van supersymmetrie zijn we aangekomen bij het doel van dit proefschrift. Nu we weten dat het Standaard Model niet de definitieve theorie is willen we aanwijzingen zoeken voor natuurkunde voorbij het Standaard Model. Samen met een grote groep fysici op CERN heb ik daarom gezocht in proton-proton botsingen naar afwijkingen van de verwachtingen van het Standaard Model, die consistent zijn met supersymmetrie. Grote afwijkingen kunnen leiden tot de ontdekking van SUSY, hetgeen een compleet nieuw beeld op de fysica in het vroege heelal zou geven, en

vele prikkelende vragen zou oproepen. Aan de andere kant leren we ook enorm veel als er geen significante afwijkingen zijn, hoewel we SUSY niet in het geheel kunnen weerleggen vanwege de vele mogelijkheden waarop het zich kan manifesteren. In dit geval kunnen we wel SUSY scenario's gedeeltelijk uitsluiten via limieten op de mogelijke massa's. Deze limieten geven de bovengrens aan op de massa's van SUSY deeltjes waarvoor SUSY uitgesloten is. Boven deze limieten kunnen we niets uitsluiten. Deze limieten bevatten veel informatie over de resterende mogelijkheden voor SUSY en zullen toekomstige fysici leiden in hun zoektocht naar nieuwe fysica.

Detecteren van deeltjes

Om te zoeken naar supersymmetrische deeltjes moet er een omgeving zijn waar ze gecreëerd en (indirect) gedetecteerd kunnen worden. Aangezien de meeste SUSY deeltjes instabiel zijn komen ze niet voor in normale materie. Om ze te kunnen ontdekken moeten ze geproduceerd worden door protonen met hoge energie op elkaar te laten botsen. Dit gebeurt in de Large Hadron Collider (LHC) op CERN in Genève. Einsteins beroemde formule $E = mc^2$ verteld ons dat energie in massa kan worden omgezet, en vice versa. Als dus twee deeltjes met hoge energie botsen kunnen een of meerdere zwaardere deeltjes gemaakt worden. De LHC is momenteel de grootste en krachtigste deeltjesversneller, met een botsingsenergie van 8 TeV, waar 1 TeV equivalent is aan de energie van een vliegende mug, bestaand uit ongeveer 10^{23} atomen.

Nadat supersymmetrische of andere zware deeltjes (zoals top quarks) zijn geproduceerd, vervallen ze onmiddellijk in lichtere quarks en leptonen. Hoewel deze deeltjes te klein zijn om direct waar te nemen kunnen ze wel indirect worden gedetecteerd. Hiervoor gebruiken we ATLAS een deeltjesdetector die om een botsingspunt van de LHC heen gebouwd is. Deeltjes die door de detector vliegen ioniseren hierin materiaal, waarna ze worden gestopt in calorimeter. Hiermee kunnen de deeltjes en hun impuls worden gereconstrueerd.

Zoeken naar supersymmetrie

Gebruikmakend van de data die op deze manier wordt verkregen uit proton-proton botsingen, of *gebeurtenissen*, heb ik gezocht naar aanwijzingen voor supersymmetrische deeltjes. Dit gebeurt via een zogenaamd 'tel-experiment'. Elke gebeurtenis wordt geclassificeerd als Standaard Model-achtig of SUSY-achtig. Het aantal waargenomen SUSY-achtige gebeurtenissen wordt vergeleken met het verwachte aantal als we uitgaan dat er enkel Standaard Model deeltjes zijn geproduceerd.

Tijdens elke botsing is er voor elk deeltje een berekenbare waarschijnlijkheid om het te creëren, afhankelijk van onder andere de massa van het deeltje. Aangezien SUSY deeltjes zwaarder zullen zijn dan hun Standaard Model partners is de kans om SUSY deeltjes te creëren kleiner dan 'normale' deeltjes. In de 20 miljoen botsingen per seconde zullen dus veel meer Standaard Model deeltjes worden gemaakt dan (hypothetische) SUSY deeltjes. De selectiecriteria voor SUSY-achtige gebeurtenissen moeten dus nauwkeurig worden geoptimaliseerd.

Selectie van de hadronische kanalen

Om onderscheid te maken tussen SUSY en Standaard Model gebeurtenissen is het essentieel om te weten wat voor signatuur deze achterlaten in de detector. In een groot gedeelte van de mogelijke SUSY scenario's zullen in de meerderheid van de gebeurtenissen twee SUSY deeltjes geproduceerd worden, die vervallen in quarks met grote impuls en twee neutralino's. Eenzame quarks zijn niet toegestaan in het Standaard Model, waardoor de uitgaande quarks hadroniseren in een cascade van deeltjes, ook wel een *jet* genoemd. De stabiele neutralino's hebben geen interactie met de detector en zijn daarom niet te detecteren. Echter, gebruikmakend van de wet van impulsbehoud kunnen we wel een netto impulsverlies in de transversale richting waarneemen: aangezien de twee inkomende protonen geen transversale impuls hebben mogen de uitgaande deeltjes ook geen netto transversale impuls hebben. Neutralino's die niet worden gedetecteerd verschijnen zo als een impuls-imbalance, die wordt gemeten als missende transversale energie $\cancel{E}_T$.

Zulke SUSY scenario's hebben aldus een *hadronische* signatuur van twee of meer jets en hoge $\cancel{E}_T$, terwijl er vaak geen leptonen worden geproduceerd. Dit is geïllustreerd in figuur S.3. Om de Standaard Model achtergronden in onze selectie te reduceren eisen we daarom dat de gebeurtenissen hoge $\cancel{E}_T$ en meerdere jets met grote impuls hebben. Aangezien in het verval van de meeste Standaard Model deeltjes wel leptonen voorkomen, reduceren we de achtergronden aanzienlijk door verder alleen gebeurtenissen zonder leptonen te selecteren. Uiteindelijk worden de selectiecriteria geoptimaliseerd voor verschillende types SUSY gebruikmakend van informatie van $\cancel{E}_T$, het aantal jets en hun transversale impuls. Dit leidt tot 12 signaal selecties, waarin de laatste eis wordt gesteld op de *effectieve massa* m_{eff} . Dit is gedefinieerd als de som van $\cancel{E}_T$ en de transversale impulsen van de waargenomen jets. De effectieve massa is gecorreleerd aan de massa van geproduceerde SUSY deeltjes en zal daarom gemiddeld hoger zijn in SUSY gebeurtenissen.

Afschatting van de achtergronden

Nadat de signaal selectiecriteria zijn vastgesteld maken we een raming van hoeveel gebeurtenissen er in de data geselecteerd zullen worden als SUSY niet bestaat. Dit aantal Standaard Model gebeurtenissen wordt geschat door een combinatie van simulaties van de data en extrapolaties van Standaard Model-rijke selecties naar onze SUSY selectie. De simulaties worden uitgevoerd door Monte Carlo generatoren, die gebeurtenissen simuleren vanuit de theoretische grondbeginselen. Ze simuleren de productie van deeltjes in een botsing, evenals hun verval en de interactie van de uitgaande deeltjes met de detector. Het is essentieel om de onzekerheden op de geraamde achtergrond te weten en te reduceren. Er zijn onzekerheden met een theoretische oorsprong, als gevolg van de simulaties, en met een experimentele oorsprong als gevolg van detector kalibraties.

Supersymmetrische deeltjes productie

Naast deze afschatting van de achtergronden moet ook het verwachte aantal SUSY signaal gebeurtenissen in de selectie worden geraamd. Hiertoe wordt ook de productie van supersymmetrische deeltjes gesimuleerd. Het aantal geproduceerde deeltjes is hierin essentieel. Zoals voor alle fundamentele deeltjes wordt de productie waarschijnlijkheid van SUSY deeltjes berekend door middel van hun *werkzame doorsnede*. De berekening van de werkzame doorsnede hangt af van theoretische en experimentele bijdragen, met als belangrijkste *parton distributie functies* (PDFs). Deze worden gebruikt om de impuls-afhankelijke waarschijnlijkheid te bepalen dat bijvoorbeeld een quark in het ene proton botst met een gluon in het andere proton. Omdat PDFs moeilijk te meten zijn hebben ze echter grote onzekerheden, die doorwerken tot onzekerheden op de berekende werkzame doorsnede. De methodes die ik heb ontwikkeld om deze vast te stellen leiden tot onzekerheden op de productiesnelheid op de productie van *squark* en *gluino* paren van tussen 10% en 70% voor deeltjes met respectievelijk lage en zeer hoge massa. Deze deeltjes zijn de SUSY partners van quarks en gluonen. Aangezien deze theoretische onzekerheden alleen invloed hebben op de SUSY signaal gebeurtenissen wordt hun effect apart getoond in de resultaten die volgen.

Resultaten

Pas wanneer al het voorgaande is voltooid kan de data, genomen in de eerste helft van 2012, getest worden tegen onze selectiecriteria. In figuur S.4 is het aantal geselecteerde gebeurtenissen uitgezet tegenover hun waargenomen m_{eff} voor een selectie met ten minste twee jets. De zwarte punten geven data aan terwijl de gekleurde gestapelde histogrammen de verschillende achtergronden weergeven. De gestreepte lijn geeft de m_{eff} verdeling weer voor een SUSY model. Als SUSY bestaat met niet te zware squarks en gluino's verwachten we grote afwijkingen te zien tussen data en achtergronden bij hoge waarden van m_{eff} , rechts van de rode pijl. Er is echter in geen enkele selectie een significante afwijking gevonden. Zonder grote afwijkingen kan het mogelijke aantal SUSY gebeurtenissen die nog in de data verscholen zijn berekend worden. Dit gebeurt met behulp van een statistisch raamwerk, gebruik makend van de onzekerheden op het geschatte aantal achtergrond gebeurtenissen. Dit kan ons vertellen welke SUSY scenario's uitgesloten kunnen worden. Zoals eerder vermeld worden zwaardere deeltjes minder vaak geproduceerd en is de verwachting dat we geen conclusies kunnen trekken voor hele hoge SUSY massa's.

Met zo veel verschillende SUSY modellen moet er een gekozen worden om onze resultaten in te interpreteren. Om (meer) model-onafhankelijk te zijn worden de uitgesloten massa's voor twee hypothetische *gesimplificeerde* SUSY modellen gegeven, waarin slechts enkele SUSY deeltjes geproduceerd kunnen worden. Links in figuur S.5 is de limiet in een model weergegeven waarin alleen squarks geproduceerd kunnen worden, terwijl rechts een model weergeeft met enkel gluino productie. Hoewel deze gesimplificeerde modellen niet de echte manifestatie van SUSY kunnen zijn, kan een combinatie van limieten in zulke modellen gebruikt worden om een limiet in elk soort SUSY model te genereren. Beide limieten zijn gegeven als functie van de squark

of gluino massa op de x -as en de neutralino massa op de y -as. De rode lijn geeft de waargenomen limiet weer: SUSY modellen met squark of gluino massa's onder deze limiet zijn uitgesloten. De gestippelde rode lijn geeft de limiet aan wanneer de werkzame doorsnede voor SUSY productie hoger of lager is, consistent met de hiervoor besproken onzekerheid. In SUSY met massaloze neutralino's moeten squarks een massa van ten minste 780 GeV hebben³, terwijl gluino's een massa van minstens 1175 GeV moeten hebben. Deze limieten zijn echter sterk afhankelijk van de neutralino massa, waarin squarks en gluino's vervallen. Voor neutralino's met een massa dichtbij die van de andere deeltjes kan SUSY niet worden uitgesloten.

Conclusie en vooruitzicht

Hoewel er helaas geen nog aanwijzingen voor supersymmetrie zijn gevonden heeft de analyse in dit proefschrift slechts een klein gedeelte van de mogelijke SUSY scenario's kunnen onderzoeken, en zijn er nog vele mogelijkheden waarin SUSY kan bestaan. Zoals hierboven vermeld is een van die mogelijkheden dat de neutralino's bijna net zo zwaar zijn als de andere SUSY deeltjes. Tegelijkertijd leidt de niet aflatende inspanning van vele natuurkundigen ertoe dat supersymmetrie langzaamaan in de hoek gedreven zal worden⁴. De inspanningen zullen in de komende jaren alleen maar toenemen nu de LHC hogere botsingsenergiën bereikt en er tot honderd keer meer data wordt verkregen. Zelfs als er geen bewijs voor SUSY wordt gevonden ben ik hoopvol dat de LHC via deze zoektochten afwijkingen in het Standaard Model zal vinden.

³Als alle partners van de vijf lichtste quarks even zwaar zijn.

⁴Althans de meest interessante vorm met niet te zware deeltjes.

Acknowledgements

My years working as a PhD has been an interesting and fun experience, but also difficult and stressful at times. This thesis would be quite a sad bunch of empty pages without the help and support I have received from so many colleagues and friends. Therefore I would like to use the last pages to thank those who helped me along the way. Since some of these are Dutch, I will switch between Dutch and English at leisure.

De eerste woorden zijn voor mijn promotor Paul. Toen ik als afgestudeerde theoreet begon te werken bij ATLAS, wist ik eigenlijk niet zo goed waar ik aan begon. Maar jij gaf me veel leesmateriaal en opdrachten die tussen de theorie en het experiment in zaten, waardoor ik langzaamaan de brug kon slaan en uiteindelijk mijn eigen weg begon te vinden binnen het experiment (waar nodig met een kleine koerswijziging door jou). Bedankt vooral ook voor je opbeurende woorden en je steun wanneer ik mijn eigen werk weer eens niet goed genoeg vond. Olya, as my co-promotor you helped me tremendously, both in the first year and (especially) during the writing of this thesis. Thank you for reading all my first, second and third drafts, with always good insight and comments.

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