Measurement of the Inclusive Electron Cross-Section from Heavy-Flavour Decays and Search for Compressed Supersymmetric Scenarios with the ATLAS Experiment

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Abstract

Subject of this thesis are the measurement of the inclusive cross-section for electron production from heavy-flavour decays and the search for compressed supersymmetric scenarios in events with missing transverse energy, jets, and one isolated low- p_T lepton in the final state. Both analyses are based on data taken by the ATLAS experiment at the Large Hadron Collider (LHC).

The measurement of the inclusive cross-section for electron production from heavy-flavour decays is motivated by the possibility to validate and constrain theoretical predictions on heavy-flavour production in the new energy regime of the LHC. In addition, the analysis allows to evaluate the performance of the ATLAS electron reconstruction in the low transverse momentum range, where the spectrum is dominated by non-isolated electrons from heavy-flavour decays. The analysis is based on the first 1.3 pb^{-1} of data collected by ATLAS in 2010 at 7 TeV center-of-mass energy in proton-proton collisions. The cross-section is measured as a function of electron transverse momentum in the range $7 \text{ GeV} < p_T < 26 \text{ GeV}$ and within $|\eta| < 2.0$ in pseudorapidity, excluding $1.37 < |\eta| < 1.52$. Theoretical predictions for heavy-flavour production obtained from Fixed Order NLO calculations with NLL resummation (FONLL) are found to be in good agreement with the measurement. Agreement is also observed with the NLO predictions of the POWHEG Monte-Carlo generator interfaced to PYTHIA for parton showering. The results are consistent with a complementary ATLAS measurement with muons in the final state. The measured value of the cross-section within the fiducial range of the analysis is

$$\sigma_{\text{HF}}^e = 0.946 \pm 0.020(\text{stat.}) \pm 0.146(\text{syst.}) \pm 0.032(\text{lumi.}) \mu\text{b}.$$

The search for compressed supersymmetric scenarios presented in the second part of this thesis is motivated by the need to explore experimentally less accessible regions of supersymmetric parameter space given the existing stringent bounds on SUSY. In compressed supersymmetric scenarios the predicted supersymmetric particles are close in mass thus leading to experimental signatures with soft physics objects in the final state. The identification of such low transverse momentum objects is exacerbated by the presence of large backgrounds in this regime. The search presented in this work is based on events with missing transverse energy, jets, and one low transverse momentum lepton in the final state. The analysis thus builds, in part, on the previously described measurement using low transverse momentum electrons. No significant excess of events over the Standard Model expectation is observed in any of the signal regions of the analysis. Model independent upper limits on the visible cross-section are calculated. Exclusion limits are derived for a number of simplified supersymmetric models with gluino and squark pair production. The sensitivity in the compressed mass region of these models is improved by up to a factor of ten with respect to two complementary ATLAS analyses with one or two high- p_T leptons in the final state. The results of all three analyses are combined statistically to provide the best exclusion limits over the full parameter range of the considered supersymmetric scenarios.

Résumé

Le travail présenté dans cette thèse est divisé en deux parties. La première est consacrée à la mesure de la section efficace inclusive de la production d'électrons provenant de la désintégration de quarks lourds. Dans la seconde partie, une recherche de modèles supersymétriques comprimés est effectuée par l'analyse d'événements avec, dans l'état final, la présence d'énergie transverse manquante, de jets et d'un lepton isolé de bas p_T . Les deux analyses sont réalisées avec des données enregistrées par l'expérience ATLAS au LHC.

La motivation pour effectuer la mesure de la section efficace inclusive de la production d'électrons provenant de la désintégration de quarks lourds est la possibilité de valider et de contraindre les prédictions théoriques pour la production de quarks lourds dans le nouveau régime d'énergie du LHC. Par ailleurs, l'analyse permet d'évaluer la performance de la reconstruction d'électrons de bas p_T du détecteur ATLAS, où le spectre est dominé par des électrons non isolés provenant de la désintégration de quarks lourds. L'analyse est réalisée avec les premières données enregistrées avec le détecteur ATLAS en 2010 avec une énergie de centre de masse de 7 TeV. La section efficace est mesurée en fonction de l'impulsion transverse des électrons dans la région $7 \text{ GeV} < p_T < 26 \text{ GeV}$ pour des valeurs de pseudorapidité de $|\eta| < 2.0$, excluant la région $1.37 < |\eta| < 1.52$. Les prédictions théoriques pour la production de quarks lourds utilisant des calculs fixes à l'ordre NLO avec resommation des termes logarithmiques à l'ordre NLL (FONLL) sont en accord avec la valeur mesurée. Les résultats sont également en conformité avec les prédictions d'ordre NLO du générateur POWHEG en combinaison avec PYTHIA pour la simulation des cascades partoniques. Les résultats sont en accord avec une analyse complémentaire utilisant des muons dans l'état final. La valeur mesurée pour la section efficace dans le régime cinématique de l'analyse est

$$\sigma_{\text{HF}}^e = 0.946 \pm 0.020(\text{stat.}) \pm 0.146(\text{syst.}) \pm 0.032(\text{lumi.}) \mu\text{b}.$$

La motivation pour la recherche de modèles supersymétriques comprimés provient de la nécessité d'explorer des régions moins accessibles expérimentalement de l'espace de paramètres supersymétriques, compte tenu des limites existantes strictes sur ces paramètres. Les modèles supersymétriques comprimés sont caractérisés par de petites différences de masse entre les particules prédites, ce qui mène à des signatures expérimentales composées d'objets avec une impulsion basse dans l'état final. L'identification de ces objets est rendue plus difficile par le large bruit du fond dans ce régime. La recherche présentée dans cette thèse est basée sur des événements avec, dans l'état final, la présence d'énergie transverse manquante, de jets et d'un lepton isolé de bas p_T . L'analyse est donc partiellement basée sur la mesure de la section efficace décrite précédemment qui utilise des électrons d'impulsion transverse basse. Aucun excès significatif d'événements n'a été observé dans les régions du signal de l'analyse. Des limites supérieures sur la section efficace visible ont été calculées indépendamment des modèles supersymétriques. Des limites d'exclusions ont été déterminées pour des modèles supersymétriques simplifiés avec la production de gluinos et squarks. La sensibilité dans les régions comprimées de ces modèles est améliorée par un facteur 10 au maximum par rapport aux deux analyses complémentaires effectuées avec l'expérience ATLAS, qui considèrent la présence d'un ou deux leptons dans l'état final. Les résultats des trois analyses sont combinés statistiquement pour obtenir la meilleure limite dans toutes les régions de l'espace de paramètres des modèles supersymétriques examinés.

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Chapter 1

Introduction

Research in the field of high energy physics explores the elementary constituents of matter and their interactions. Over the past decades, enormous experimental and theoretical efforts have led to a consistent and very accurate theoretical description of the physical observations - the Standard Model of particle physics. Therein, twelve point-like fermions and their corresponding antiparticles constitute the basic building blocks of matter. The interactions between them are mediated by four types of gauge bosons (ignoring charges) which manifest themselves in the weak, electromagnetic, and strong forces.

Within the Standard Model, the elementary particles acquire their mass through the mechanism of electroweak symmetry breaking which predicts the existence of an additional particle - the Higgs boson. The tremendous experimental efforts in the search for this last missing piece of the Standard Model have recently culminated in the discovery of a new particle which, pending further scrutiny, could indeed be the predicted boson.

Despite the outstanding success of the Standard Model, it is not expected to be the ultimate theory of particle physics. Reasons for this are, amongst many others, the failure of the Standard Model's formalism to incorporate the gravitational interaction, the lack of an explanation for dark matter which makes up a large fraction of our universe, the non-unification of the electroweak and strong coupling constants when extrapolated to high energies, and the fine-tuning of radiative corrections to the mass of the Higgs boson. One candidate theory, that provides convenient solutions to many of the open questions of the Standard Model, is Supersymmetry (SUSY). SUSY is a fundamental symmetry between fermions and bosons, that predicts a set of new partner particles to the known Standard Model particles. While no supersymmetric particles have yet been observed by experiment, many theoretically motivated arguments favour their existence at the electroweak energy scale.

With the start-up of the Large Hadron Collider (LHC) in 2010 at the European Organization for Nuclear Research (CERN), a new energy regime has become accessible. During the first three years of operation at the energy frontier, the LHC has delivered approximately 60 fb^{-1} of proton-proton collisions at 7 and 8 TeV center-of-mass energy. These particle collisions are recorded and analyzed by large-scale experiments built around the interaction points of the LHC. One of these experiments is ATLAS, a general-purpose detector designed to exploit the full physics potential of the LHC.

The performance of ATLAS in the search for new physics scenarios depends critically on the

understanding of the experimental apparatus and the accuracy of the theoretical predictions for the expected Standard Model background processes. Particularly during the commissioning phase of ATLAS, it was therefore important to perform precision measurements of the known Standard Model processes to calibrate the detector and to validate the theoretical predictions in this new energy regime.

Part of this effort was the first physics analysis presented in this thesis, the measurement of the inclusive cross-section for electron production from decays of heavy-flavour bottom and charm quarks using 1.3 pb^{-1} of data collected with ATLAS in 2010 at 7 TeV center-of-mass energy. As one of the first analyses performed in ATLAS, this measurement constitutes not only a powerful test of heavy-flavour QCD predictions at LHC energies, but also promotes the general understanding of the electron reconstruction performance of ATLAS in the low transverse momentum range. The latter aspect is of particular relevance in the context of new physics searches where the discovery potential can be significantly increased if leptons with low transverse momentum are considered.

An example for such an analysis is the search for compressed supersymmetric scenarios presented in the second part of this thesis. In compressed SUSY the mass differences between the predicted particles are small and the transverse momentum spectra of the observable decay products are soft. The search is therefore based on the presence of one isolated low transverse momentum lepton in addition to jets and missing transverse energy in the final state. The data used in the analysis amounts to 4.7 fb^{-1} collected with ATLAS in 2011 at 7 TeV center-of-mass energy. The motivation to explore the experimentally less accessible compressed regime of supersymmetric parameter space arises from the existing stringent bounds on SUSY from earlier LHC searches.

This thesis is organized as follows. Chapter 2 describes the formalism of the Standard Model of particle physics and gives an introduction to the theory of Supersymmetry. Chapter 3 outlines the main features of the LHC and the ATLAS detector which form the experimental setup of this thesis. The acquisition and processing of data by ATLAS, including the reconstruction of physics object and quantities, and the generation of simulated event samples, are reviewed in chapter 4. This chapter also includes a description of the Trigger-Tier0 liason task that the author fulfilled during the first two ATLAS data-taking campaigns. Chapter 5 describes the first physics analysis carried out within this thesis, the measurement of the inclusive cross-section for electron production from heavy flavour decays using 2010 ATLAS data. Chapter 6 then reports on the search for compressed supersymmetric scenarios in final states with missing transverse energy, jets, and one isolated lepton with low transverse momentum using the full 2011 ATLAS dataset. Finally, the results of this thesis are summarized in Chapter 7.

Chapter 2

Theoretical Background

This chapter reviews the theoretical foundation of the work presented in this thesis. Section 2.1 outlines the main features of the Standard Model of particle physics. Section 2.2 gives an introduction to Supersymmetry. The content is taken from the referenced sources. Parts of this chapter are taken from the diploma thesis of the author [1] and have been adapted according to the latest developments.

2.1 The Standard Model of Particle Physics

Today’s experimentally verified knowledge of the fundamental particles and their interactions is summarized in the Standard Model of particle physics (SM).

Within the Standard Model, the elementary constituents of matter are twelve spin-1/2 fermions and their respective antiparticles. They can be further classified according to their interactions into quarks and leptons, for each of which three generations of particle pairs exist. For both fermion species the second and third generations are heavier¹ copies of the first generation with identical quantum numbers. A summary of the Standard Model fermions is shown in Table 2.1 (left).

Quarks participate in the strong, weak and electromagnetic interactions and therefore carry color, electric charge and weak isospin. For each generation there is one up-type quark with electric charge +2/3 (“up”, “charm”, “top”) and one down-type quark with -1/3 (“down”, “strange”, “bottom”). The lepton generations consist of one electron-type lepton (“electron”, “muon”, “tau”) with electric charge -1 and a neutral almost massless lepton-neutrino. Whereas electron-type leptons are charged and participate in electromagnetic and weak interactions, neutrinos can only interact weakly due to their lack of electric charge.

The interactions between the matter particles are mediated by the spin-1 gauge bosons summarized in Table 2.1 (right). Their existence arises from invariance of the respective interaction Lagrangians under local symmetry transformations. The underlying mathematical structure is the direct product of the internal symmetry groups $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$, where $SU(3)_C$ describes the strong force mediated by eight gluons and $SU(2)_L \otimes U(1)_Y$ the unified electroweak force mediated by $W^\pm$, Z bosons and the photon γ . Gravitational interactions are not incor-

¹Not yet established for neutrinos.

Fermions	1. Generation	2. Generation	3. Generation	Bosons
Quarks	u $m=2.3^{+0.7}_{-0.5}$ MeV $Q=2/3$	c $m=1.28\pm0.03$ GeV $Q=2/3$	t $m=173.5\pm0.6\pm0.8$ GeV $Q=2/3$	γ $m=0$ $Q=0$
	d $m=4.8^{+0.7}_{-0.3}$ MeV $Q=-1/3$	s $m=95^{+5}_{-5}$ MeV $Q=-1/3$	b $m=4.19\pm0.03$ GeV $Q=-1/3$	g $m=0$ $Q=0$
Leptons	ν_e $m<2.05$ eV $Q=0$	ν_μ $m<0.19$ MeV $Q=0$	ν_τ $m<18.2$ MeV $Q=0$	Z^0 $m=91.19\pm0.002$ GeV $Q=0$
	e $m=0.511$ MeV $Q=-1$	μ $m=106$ MeV $Q=-1$	τ $m=1.78$ GeV $Q=-1$	$W^\pm$ $m=80.39\pm0.02$ GeV $Q=\pm1$

Table 2.1: Summary of the experimentally measured spin- $\frac{1}{2}$ fermions (left) and the spin-1 gauge bosons (right) of the Standard Model of particle physics. The numbers in small font correspond to the electric charge Q in units of the elementary charge e and the mass m of the particles as given by [2]. The very small uncertainties on the masses of the charged leptons are not quoted. The existence of the spin-0 Standard Model Higgs boson has not yet been fully established at the time of writing and is therefore not listed here.

porated into the Standard Model.

The unification of the weak and electromagnetic forces happens by virtue of the so-called Glashow-Salam-Weinberg (GSW) mechanism [3–5], in which a local gauge transformation of the left-handed weak isospin doublets $SU(2)_L$ and the $U(1)_Y$ multiplets and singlets with respect to hypercharge Y takes place. The invariance of the Lagrangian under these transformations necessitates the introduction of new massless vector fields $W^{1,2,3}$ and B , of which the physically observed $W^\pm$, Z and γ bosons are linear combinations. The physical eigenstates of the latter two bosons are additionally rotated by a so-called Weinberg-angle θ_W .

The Standard Model particles acquire their masses through interaction with a scalar Higgs background field, which spontaneously breaks the electroweak symmetry. The introduction of this breaking mechanism is necessary since gauge invariance requires massless gauge bosons, which contradicts experimental results. The neutral spin-0 Higgs boson associated to this mechanism has been the subject of many experimental searches. Recently the discovery of a new particle has been reported which could finally bring the long awaited experimental verification of this last missing piece of the Standard Model. Further details on this subject are given in section 2.1.6.

The mathematical framework behind the Standard Model is a relativistic quantum field theory which can be derived using the Lagrangian formalism. The starting point is the Lorentz invariant scalar Lagrangian density $\mathcal{L}$ which describes the dynamics of the system of interest. The application of Hamilton’s principle of least action, $\delta S = 0$, where the action S is defined as

$$S = \int d^4x \mathcal{L}, \quad (2.1)$$

then leads to the Euler-Lagrange equations from which the explicit equations of motion for the considered fields can be derived. While every global symmetry of the Lagrangian up to a total derivative leads to a conservation law according to Noether’s theorem [6], invariance with respect to local gauge transformations implies the introduction of new gauge fields which

mediate the fundamental interactions of the Standard Model.

In the following sections the theory of the electromagnetic, weak, and strong interactions of the Standard Model as well as the Higgs mechanism is derived using the Lagrangian formalism.

2.1.1 Electromagnetic Interaction

The free propagation of the fermions of the Standard Model is described by the Dirac equation,

$$(i\gamma^\mu \partial_\mu - m)\psi = 0. \quad (2.2)$$

Here ψ is the four-component spinor representing the fermionic field, m the rest mass of the associated particle, and γ^μ are the Dirac matrices, defined as

$$\gamma^0 = \begin{bmatrix} \mathbb{1}_2 & 0 \\ 0 & \mathbb{1}_2 \end{bmatrix}, \quad \gamma^k = \begin{bmatrix} 0 & \sigma^k \\ -\sigma^k & 0 \end{bmatrix}, \quad k = 1, 2, 3, \quad (2.3)$$

where σ^k are the Pauli matrices,

$$\sigma^1 = \begin{bmatrix} 0 & 1 \\ 1 & 0 \end{bmatrix}, \quad \sigma^2 = \begin{bmatrix} 0 & -i \\ i & 0 \end{bmatrix}, \quad \sigma^3 = \begin{bmatrix} 1 & 0 \\ 0 & -1 \end{bmatrix}. \quad (2.4)$$

Every component of the spinor ψ is a function of the space-time coordinates, which in their contravariant form are written as

$$x^\mu = (t, x, y, z) = (x^0, x^1, x^2, x^3), \quad (2.5)$$

with the corresponding covariant derivative defined as

$$\partial_\mu = \frac{\partial}{\partial x^\mu} = \left(\frac{\partial}{\partial t}, \vec{\nabla} \right) = \left(\frac{\partial}{\partial x^0}, \frac{\partial}{\partial x^1}, \frac{\partial}{\partial x^2}, \frac{\partial}{\partial x^3} \right). \quad (2.6)$$

In the Lagrangian formalism the Dirac equation 2.2 follows from the Lagrangian

$$\mathcal{L} = \bar{\psi}(i\gamma^\mu \partial_\mu - m)\psi, \quad (2.7)$$

with the adjoint spinor $\bar{\psi}$ defined as

$$\bar{\psi} = \psi^\dagger \gamma^0, \quad (2.8)$$

where $\psi^\dagger$ is the hermitian conjugate of the spinor. This Lagrangian is invariant under a global transformation of the type

$$\psi \rightarrow \psi' = e^{iq\alpha} \psi \quad (2.9)$$

which, following Noether's theorem, leads to the conservation of the electromagnetic current

$$j^\mu = q\bar{\psi}\gamma^\mu\psi \quad (2.10)$$

with the electric charge q . Here q is defined as $q = Qe$, where Q is the charge quantum number of the particle involved (see table 2.1) and e is the elementary charge. The transformations of equation 2.9 with the real parameter α build the group of unitary transformations $U(1)$.

To introduce the electromagnetic interaction to the theory the symmetry of the Lagrangian has to be extended to local gauge transformations of the type

$$\psi \rightarrow \psi' = e^{iq\alpha(x)}\psi, \quad (2.11)$$

where the parameter α is now dependent on the position in space-time. To preserve the invariance of the Lagrangian density under such transformations, the derivative ∂_μ in equation 2.7 must be replaced by a new covariant derivative

$$D_\mu = \partial_\mu + iqA_\mu. \quad (2.12)$$

This expression contains a newly introduced vector field A_μ which is required to transform as

$$A_\mu \rightarrow A_\mu - \partial_\mu\alpha(x). \quad (2.13)$$

With these modifications the Lagrangian density of equation 2.7 reads

$$\mathcal{L} = \bar{\psi}(i\gamma^\mu D_\mu - m)\psi = \bar{\psi}i\gamma^\mu\partial_\mu\psi - m\bar{\psi}\psi - q\bar{\psi}\gamma^\mu\psi A_\mu. \quad (2.14)$$

Here the first two terms are equivalent to equation 2.7 and describe the free propagation of the fermion fields of the Standard Model. The last term accounts for the newly found interaction of the charged fermions with the vector field A_μ which can be identified with the photon γ . The coupling strength of the interaction is given by the electric charge q of the fermion. Since a mass term of the form $\mathcal{L}_\gamma = \frac{1}{2}m^2 A^\mu A_\mu$ is not gauge invariant, the theory requires the photon to be massless in accordance with experimental observation.

To complete the Lagrangian of the electromagnetic interaction a gauge invariant kinetic term for A_μ needs to be added to equation 2.14. The kinematics of the photon field are governed by Maxwell's equations, which in their covariant form can be expressed in terms of the field strength tensor

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu. \quad (2.15)$$

The kinetic term of the photon in the Lagrangian, from which Maxwell's equations can be derived, is given by

$$\mathcal{L} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} \quad (2.16)$$

The final Lagrangian density of quantum electrodynamics is thus given by:

$$\mathcal{L}_{\text{QED}} = \underbrace{\bar{\psi}i\gamma^\mu\partial_\mu\psi}_{\text{fermion kinetic term}} - \underbrace{m\bar{\psi}\psi}_{\text{fermion mass term}} - \underbrace{\frac{1}{4}F^{\mu\nu}F_{\mu\nu}}_{\text{photon kinetic term}} - \underbrace{q\bar{\psi}\gamma^\mu\psi A_\mu}_{\text{interaction term}} \quad (2.17)$$

2.1.2 Weak Interaction and Electroweak Unification

The first theory of the weak interaction was proposed by Fermi in 1933 [7]. The theory was motivated by the earlier postulation of the neutrino in 1927 by Pauli to explain the continuous

energy spectrum of electrons from β decays. The Fermi theory describes the weak interaction with four-fermion vertices which is a valid approximation at energies much lower than the mass of the W-boson. The Lagrangian of this interaction is of the type current-current

$$\mathcal{L}_{\text{Fermi}} = \frac{G_F}{\sqrt{2}} J_\mu^\dagger J^\mu, \quad (2.18)$$

where the current is given by

$$J^\mu = \bar{\psi} \gamma^\mu \psi, \quad (2.19)$$

and $\psi, \bar{\psi}$ are spinors associated with the fermions of the interaction. The value of the constant G_F can be obtained for example from measurements of the muon lifetime [8, 9]. However, the form of the Lagrangian of equation 2.18 is not complete since it does not account for the parity violating nature of the weak interactions, which was first suspected by Lee and Yang [10] and then experimentally confirmed by Wu et al. in 1956 [11]. The Wu experiment observed that in β -decays of Cobalt 60 nuclei the emission of electrons occurs preferentially opposite to the direction of the spin of the nucleus. Since the electron momentum transforms like a vector and the spin like an axial vector this means that parity is not conserved in this process. This effect is taken care of in the so-called V-A (vector minus axial vector) extension [12, 13] of Fermi's theory, which introduces chirality operators

$$\psi_L = P_L \psi = \frac{1}{2}(1 - \gamma^5)\psi \quad \psi_R = P_R \psi = \frac{1}{2}(1 + \gamma^5)\psi, \quad (2.20)$$

which project out the components of left- and right-handed chirality of the spinor ψ respectively. The matrix γ^5 is defined as

$$\gamma^5 = i\gamma^0\gamma^1\gamma^2\gamma^3. \quad (2.21)$$

and the weak current is then of the type

$$J^\mu = \bar{\psi}_e \gamma^\mu \frac{(1 - \gamma^5)}{2} \psi_\nu, \quad (2.22)$$

where the example of the leptonic $e - \nu_e$ process has been chosen. Equation 2.22 has the form vector current minus axial-vector current which explains the name V-A theory. The simultaneous appearance of the two types of currents means that parity is not conserved in weak processes and only the left-chiral component of the fermions and the right-chiral component of anti-fermions participate in the interaction.

As mentioned above the Fermi theory as well as its V-A extension is only valid in a low energy regime. Above a certain threshold the calculated cross-section of a given process rises quickly with its energy which violates the unitarity of the theory.

The solution to this problem is the introduction of intermediate vector bosons, which, as in the electromagnetic interactions, arise from invariance of the Lagrangian under a given set of local gauge transformations. In fact, as will be shown in the following, a complete and renormalizable

theory of the weak and electromagnetic interactions requires their unification into a common electroweak framework [3–5] based on the symmetry groups $SU(2)_L \otimes U(1)_Y$.

To describe the weak interactions with parity violation the symmetry $SU(2)_L$ is assumed. The left-handed fermions are ordered in doublets of weak isospin, e.g. for the first generation of leptons

$$\psi_L = \begin{pmatrix} \psi_{1,L} \\ \psi_{2,L} \end{pmatrix} = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix}. \quad (2.23)$$

In analogy to equation 2.7 for the electromagnetic interactions, the Lagrangian density of two free, massless, left-handed fermions is given by

$$\mathcal{L}_L = \bar{\psi}_{1,L}(i\gamma^\mu\partial_\mu)\psi_{1,L} + \bar{\psi}_{2,L}(i\gamma^\mu\partial_\mu)\psi_{2,L} = \bar{\psi}_L(i\gamma^\mu\partial_\mu)\psi_L. \quad (2.24)$$

To obtain the mediators of the weak interaction, invariance of the Lagrangian under local gauge transformations of the form

$$\psi_L \rightarrow \psi'_L = e^{ig\alpha^a(x)T^a}\psi_L \quad (2.25)$$

is required. The 2×2 matrices T^a are the generators of $SU(2)_L$ and can be chosen as

$$T^a = \frac{\sigma^a}{2}, \quad a = 1, 2, 3, \quad (2.26)$$

where σ^a are the Pauli matrices given in equation 2.4. The $\alpha(x)$ are arbitrary space and time dependent functions, and g will be interpreted as the coupling constant of the interaction.

To preserve invariance of the Lagrangian under $SU(2)_L$ transformations, a new covariant derivative

$$D_\mu = \partial_\mu + igT^aW_\mu^a \quad (2.27)$$

is introduced along with three new vector boson fields W_μ^a , which are required to transform according to

$$W_\mu^a \rightarrow W_{\mu'}^a = W_\mu^a + \partial_\mu\alpha^a(x) - g\epsilon^{abc}\alpha^b(x)W_\mu^c, \quad (2.28)$$

where ϵ is the Levi-Civita tensor. To account for the kinetic energy of the newly found gauge fields a tensor of the form

$$W_{\mu\nu}^a = \partial_\mu W_\nu^a - \partial_\nu W_\mu^a - g\epsilon^{abc}W_\mu^bW_\nu^c \quad (2.29)$$

is defined. The complete, gauge invariant Lagrangian of $SU(2)_L$ is then written as

$$\begin{aligned} \mathcal{L}_L &= \bar{\psi}_L(i\gamma^\mu D_\mu)\psi_L - \frac{1}{4}W^{a\mu\nu}W_{\mu\nu}^a \\ &= \underbrace{\bar{\psi}_L(i\gamma^\mu\partial_\mu)\psi_L}_{\text{fermion kinetic term}} - \underbrace{g\bar{\psi}_L(\gamma^\mu T^a W_\mu^a)\psi_L}_{\text{fermion - vector field interaction term}} - \underbrace{\frac{1}{4}W^{a\mu\nu}W_{\mu\nu}^a}_{\text{vector field kinetic and self-interaction term}} \end{aligned} \quad (2.30)$$

The first term describes the kinetic energy of the left-handed fermions, the second term their interaction with the vector fields, and the third term the kinetic energy of the vector fields and their self-interaction. The self-interaction of the vector fields originates from the fact that, contrary to $U(1)$ in the case of the electromagnetic interactions, the symmetry group $SU(2)_L$ is non-abelian i.e. its generators (equation 2.26) do not commute.

While the $SU(2)_L$ gauge theory alone describes many aspects of the weak interactions it is not a complete theory. For instance it is not able to explain the experimentally observed masses of the fermions and gauge bosons associated with the vector fields and it lacks a consistent description of the interactions mediated by the neutral Z boson.

To resolve these problems the electroweak formalism introduces a combination of $SU(2)_L$ and a new $U(1)_Y$ symmetry group, where Y is the so-called hypercharge. The behaviour of the Lagrangian under $U(1)_Y$ gauge transformations is identical to that of the electromagnetic interactions and leads to a covariant derivative of the form

$$D_\mu = \partial_\mu + ig' \frac{Y}{2} B_\mu, \quad (2.31)$$

with a new vector field B_μ and a new coupling constant g' . In analogy to equation 2.15 the field strength tensor

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu \quad (2.32)$$

is defined. In contrast to the $SU(2)_L$, the $U(1)_Y$ symmetry applies not only to particles of the left-handed doublets (eq. 2.23) but also to right-handed particles which are ordered in isospin singlets. The $U(1)_Y$ invariant Lagrangian can thus be written as:

$$\mathcal{L}_Y = \bar{\psi}(i\gamma^\mu \partial_\mu)\psi + g' \frac{Y_R}{2} \bar{\psi}_R(\gamma^\mu B_\mu)\psi_R + g' \frac{Y_L}{2} \bar{\psi}_L(\gamma^\mu B_\mu)\psi_L - \frac{1}{4} B^{\mu\nu} B_{\mu\nu}, \quad (2.33)$$

where the couplings depend on the type and chirality of the particle considered.

The combined electroweak Lagrangian of the $SU(2)_L$ and $U(1)_Y$ gauge theories is then simply obtained from the sum of the Lagrangians of equations 2.30 and 2.33. With the combined covariant derivative

$$D_\mu = \partial_\mu + igT^a W_\mu^a + ig' \frac{Y}{2} B_\mu, \quad (2.34)$$

written out the complete $\mathcal{L}_{\text{electroweak}}$ is given by:

$$\mathcal{L}_{\text{electroweak}} = \underbrace{-\frac{1}{4} W^{a\mu\nu} W_{\mu\nu}^a - \frac{1}{4} B^{\mu\nu} B_{\mu\nu}}_{\text{gauge field kinetic energy and self-interaction terms}} + \underbrace{\bar{\psi}_L \gamma^\mu (i\partial_\mu - \frac{1}{2} g\sigma^a W_\mu^a - \frac{1}{2} g' Y B_\mu) \psi_L}_{\text{left-chiral fermion kinetic energy and interaction term}} + \underbrace{\bar{\psi}_R \gamma^\mu (i\partial_\mu - \frac{1}{2} g' Y B_\mu) \psi_R}_{\text{right-chiral fermion kinetic energy and interaction term}} \quad (2.35)$$

To relate this Lagrangian to the physically observed charged and neutral currents of the electroweak interactions the $SU(2)_L$ generators (eq. 2.26) and thus the Pauli matrices are inserted into the covariant derivative 2.34 which then becomes

$$\begin{aligned} D_\mu &= \partial_\mu + \frac{ig}{2} \left(\begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix} W_\mu^1 + \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix} W_\mu^2 + \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix} W_\mu^3 \right) + ig' \frac{Y}{2} B_\mu \\ &= \partial_\mu + \frac{ig}{2} \begin{pmatrix} W^3 & W^1 - iW^2 \\ W^1 + iW^2 & -W^3 \end{pmatrix}_\mu + ig' \frac{Y}{2} \begin{pmatrix} B & 0 \\ 0 & B \end{pmatrix}_\mu. \end{aligned} \quad (2.36)$$

The expression is simplified by introducing the ladder operators

$$T^\pm = \frac{1}{\sqrt{2}}(T^1 \pm iT^2) \quad \text{and} \quad W^\pm = \frac{1}{\sqrt{2}}(W^1 \mp iW^2), \quad (2.37)$$

where the $W^\pm$ are identified with the charged gauge bosons. Equation 2.36 can then be written as

$$\begin{aligned} D_\mu &= \partial_\mu + \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & W^+ \\ W^- & 0 \end{pmatrix}_\mu + \frac{1}{2} \begin{pmatrix} gW^3 + g'YB & 0 \\ 0 & -gW^3 + g'YB \end{pmatrix}_\mu \\ &= \partial_\mu + \underbrace{(T^+W^+ + T^-W^-)_\mu}_{D_\mu^W} + \underbrace{(gT^3W^3 + g'\frac{Y}{2}B)_\mu}_{D_\mu^Z}, \end{aligned} \quad (2.38)$$

where the first term inserted into the Lagrangian 2.35 describes the interactions of the $W^\pm$ gauge bosons with fermions

$$\begin{aligned} \mathcal{L}_{L, \text{ charged}} &= ig\bar{\psi}_L i\gamma^\mu (T^+W^+ + T^-W^-)_\mu \psi_L \\ &= \frac{1}{\sqrt{2}} ig(\bar{\nu}_e, \bar{e})_L i\gamma^\mu \begin{pmatrix} 0 & W^+ \\ W^- & 0 \end{pmatrix}^\mu \begin{pmatrix} \nu_e \\ e \end{pmatrix}_L \\ &= -\frac{1}{\sqrt{2}} g(\bar{\nu}_{e,L} \gamma^\mu W_\mu^+ e_L + \bar{e}_L \gamma^\mu W_\mu^- \nu_{e,L}). \end{aligned} \quad (2.39)$$

The interaction of the fermions with the neutral gauge bosons originates from the term D_μ^Z , which leads to the neutral interaction Lagrangian of the form

$$\begin{aligned} \mathcal{L}_{L, \text{ neutral}} &= \bar{\psi}_L i\gamma^\mu i(gT^3W_\mu + g'\frac{Y}{2}B_\mu) \psi_L + \bar{\psi}_R i\gamma^\mu i(g'\frac{Y}{2}B_\mu) \psi_R \\ &= \sum_{\psi=e_L, e_R, \nu_L, \nu_R} \bar{\psi} i\gamma^\mu i(gT^3W_\mu + g'\frac{Y}{2}B_\mu) \psi. \end{aligned} \quad (2.40)$$

To obtain the physically observed photon and Z boson a change of base can be performed of the type

$$\begin{pmatrix} W_\mu^3 \\ B_\mu \end{pmatrix} = \begin{pmatrix} \cos \theta_W & \sin \theta_W \\ -\sin \theta_W & \cos \theta_W \end{pmatrix} \begin{pmatrix} Z_\mu \\ A_\mu \end{pmatrix}. \quad (2.41)$$

This corresponds to a rotation by the so-called Weinberg angle θ_W with which one obtains the interaction Lagrangian for the neutral current

$$\begin{aligned}
\mathcal{L}_{\gamma,Z} = & \sum_{e\nu} \bar{\psi} i \gamma^\mu i \underbrace{(g \sin \theta_W T^3 + g' \frac{Y}{2} \cos \theta_W)}_{\text{fermion-photon coupling}} A_\mu \psi \\
& + \sum_{e\nu} \bar{\psi} i \gamma^\mu i \underbrace{(g \cos \theta_W T^3 + g' \frac{Y}{2} \sin \theta_W)}_{\text{fermion-Z coupling}} Z_\mu \psi
\end{aligned} \tag{2.42}$$

Identifying the field A_μ with the photon with coupling constant equal to the Q times electric charge e leads to the expression

$$Qe = g \sin \theta_W T^3 + g' \frac{Y}{2} \cos \theta_W \tag{2.43}$$

Together with the Gell-Mann Nishijima equation [14, 15]

$$Q = T^3 + \frac{Y}{2}. \tag{2.44}$$

one obtains the relation between the Weinberg angle and the coupling constants:

$$e = g \sin \theta_W = g' \cos \theta_W \tag{2.45}$$

For the coupling of the Z boson to the fermions one obtains:

$$g_Z = g \cos \theta_W T^3 + g' \frac{Y}{2} \sin \theta_W = \frac{e}{\sin \theta_W \cos \theta_W} (T^3 - \sin^2 \theta_W Q) \tag{2.46}$$

With the above relations the interaction Lagrangian (eq. 2.42) can be rewritten as

$$\mathcal{L}_{\gamma,Z} = \sum_{e\nu} i \bar{\psi} i \gamma^\mu (Qe A_\mu + g_Z Z_\mu) \psi \tag{2.47}$$

Since the field Z_μ is a combination of the field W_μ^3 and B_μ it couples to left- and right-handed fermions.

2.1.3 Spontaneous Symmetry Breaking and Higgs Mechanism

The formalism developed in the previous sections describes all phenomena of the electroweak interactions except for the masses of the gauge bosons and fermions. The particle masses can be incorporated into the Standard Model by introducing the Higgs mechanism [16–21] which spontaneously breaks the electroweak symmetry.

The Higgs-mechanism postulates a new complex scalar $SU(2)_L$ doublet field with hypercharge $Y = 1$. The electric charge Q of the two field components follows from the Gell-Mann Nishijima relation 2.44:

$$\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} \phi_1 + i\phi_2 \\ \phi_3 + i\phi_4 \end{pmatrix} \tag{2.48}$$

Figure 2.1: *Visualization of the Higgs potential in the complex (ϕ_1, ϕ_2) -plane. The green circle indicates the position of the degenerate minima of the potential.*

As can be seen above the complex nature of the fields leads to 4 degrees of freedom. The Lagrangian density of the spin-0 doublet follows the Klein-Gordon equation

$$\mathcal{L}_\phi = |D^\mu \phi|^2 - V(\phi) \equiv (D^\mu \phi)^\dagger (D_\mu \phi) - V(\phi), \quad (2.49)$$

where D_μ is the previously introduced covariant derivative of the electroweak interactions

$$D_\mu = \partial_\mu + igT^a W_\mu^a + ig' \frac{Y}{2} B_\mu, \quad (2.50)$$

and $V(\phi)$ is a newly postulated potential which is required to be invariant under local gauge transformations. This requires the potential to be symmetric in all four components which leads to the ansatz

$$V(\phi) = -\mu^2 |\phi|^2 + \lambda |\phi|^4, \quad (2.51)$$

where μ^2 and λ are real constants. For positive values of μ^2 the potential has the characteristic “Mexican Hat” type shape shown in Fig. 2.1. The potential has a circle of degenerate minima in the complex (ϕ_1, ϕ_2) -plane with a radius of

$$v = \sqrt{\frac{\mu^2}{\lambda}}, \quad (2.52)$$

which is referred to as the vacuum expectation value of the Higgs field.

The choice of a particular ground state of the vacuum on this circle breaks the symmetry spontaneously. Since the vacuum is electrically neutral the ground state

$$\langle \phi^+ \rangle = 0, \quad \langle \phi^0 \rangle = v, \quad (2.53)$$

can be chosen. An expansion around this ground state gives

$$\phi(x) = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 + \eta_+(x) + i\xi_+(x) \\ v + \eta_0(x) + i\xi_0(x) \end{pmatrix}. \quad (2.54)$$

The fields $\eta_+(x)$, $\xi_+(x)$, $\xi_0(x)$ lead to massless Goldstone bosons [22, 23] which can be absorbed with a local gauge transformation of the type

$$\phi \rightarrow \phi' = e^{ig\alpha^a(x)T^a} \phi, \quad (2.55)$$

where the arbitrary functions $\alpha(x)$ are chosen accordingly. This gauge transformation results in the three spin degrees of freedom of the vector bosons of electroweak theory. The field ϕ is now given by

$$\phi(x) = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + \eta^0(x) \end{pmatrix}, \quad (2.56)$$

where $\eta_0(x)$ can be associated with the field of the Higgs particle and will be denoted H in the following. Substituting equation 2.56 into the potential given in equation 2.51 and using equation 2.53 leads to

$$V(\phi) = \frac{1}{2}(2\mu^2)H^2 + \lambda v H^3 + \frac{1}{4}\lambda H^4, \quad (2.57)$$

where the first term describes the mass of the Higgs boson

$$m_H = \sqrt{2\mu^2}, \quad (2.58)$$

and the second and third terms are associated with the self-interactions of the Higgs-field with 3 and 4 particle vertices.

The interaction of the Higgs boson with the vector boson of the electroweak theory is determined from the covariant derivative term of the Lagrangian

$$\begin{aligned} \mathcal{L}_\phi &= |D_\mu \phi|^2 - V(\phi) \\ &= |(\partial_\mu + D_\mu^W + D_\mu^{\gamma Z})\phi|^2 - V(\phi) \\ &= |(\partial_\mu + ig(T^+ W^+ + T^- W^-)_\mu + iQeA_\mu + ig_Z Z_\mu)\phi|^2 - V(\phi) \end{aligned} \quad (2.59)$$

Due to the photon being massless all terms including the photon field A^μ are evaluated to zero. Also no mixing terms between $W^\pm$ and Z fields occurs since $\phi^+ = 0$ was chosen. Expression 2.59 thus becomes

$$\begin{aligned}\mathcal{L}_\phi &= |\partial_\mu|^2 + |D_\mu^W|^2 + |D_\mu^{\gamma Z}\phi|^2 - V(\phi) \\ &= \frac{1}{2}(\partial^\mu H)(\partial_\mu H) + \frac{1}{4}g^2 W_\mu^+ W_\mu^- (v + H)^2 + \frac{1}{2}g_Z^2 Z^\mu Z_\mu (v + H)^2 - V(\phi)\end{aligned}\quad (2.60)$$

The first term represents the kinetic energy of the Higgs field H . Since g and v are constants, all terms with v^2 can be interpreted as mass terms, where the mass values are given by

$$m_W = \frac{1}{2}gv \quad \text{and} \quad m_Z = \frac{1}{2}\sqrt{g^2 + g'^2}v \quad (2.61)$$

for the W and Z bosons respectively. The terms

$$\frac{1}{2}g^2 W^+ W^- H \quad \text{and} \quad \frac{1}{4}g_Z^2 Z Z H \quad (2.62)$$

demonstrate that the Higgs - vector boson couplings are proportional to the masses of the vector bosons. Finally the terms

$$\frac{1}{4}g^2 v W^+ W^- H H \quad \text{and} \quad \frac{1}{2}g_Z^2 v Z Z H H \quad (2.63)$$

correspond to the interactions of two Higgs bosons and two vector bosons.

Since the value of μ in equation 2.58 is not known, the mass of the Higgs cannot be predicted from theory. The value of v can be derived e.g. from the experimentally measured values of M_Z and g_Z

$$v = \frac{2M_Z}{g_Z} = 246 \text{ GeV}. \quad (2.64)$$

While the formalism developed above incorporates the masses of the electroweak gauge bosons into the theory, the masses of the Standard Model fermions cannot be explained in the same way. In addition, standard fermion mass terms of the type

$$\mathcal{L}_m = -m\psi\psi = -m(\bar{\psi}_R\psi_L + \bar{\psi}_L\psi_R) \quad (2.65)$$

are not invariant under local $SU(2)_L$ or $U(1)_Y$ gauge transformations. This leads to the formulation of the so-called Yukawa interaction, which in the case of leptons it is given by

$$\mathcal{L}_{\text{Yukawa}} = -\lambda_l(\bar{L}_L\phi^0 L_R + \bar{L}_R\phi^+ L_L), \quad (2.66)$$

where L_L are the left-handed lepton doublets and L_R the lepton singlets. The Yukawa term is automatically invariant under $SU(2)_L$ since both ϕ and L belong to the same $SU(2)_L$ doublet. Looking at the first generation of leptons and using expression 2.56 one obtains

$$\mathcal{L}_{\text{Yukawa}} = -\lambda_e \frac{v}{\sqrt{2}} (\bar{e}_R e_L + \bar{e}_L e_R) - \lambda_e \frac{H}{\sqrt{2}} (\bar{e}_R e_L + \bar{e}_L e_R), \quad (2.67)$$

containing the electron mass term with mass

$$m_e = \lambda_e \frac{v}{\sqrt{2}} \quad (2.68)$$

and the coupling to the Higgs field proportional to that mass. Similar expressions can be derived for the remaining Standard Model fermions. In the case of the quarks the so-called Cabibbo-Kobayashi-Maskawa (CKM) matrix [24, 25] has to be introduced to take care of the fact that the mass eigenstates are rotated with respect to the flavour eigenstates.

In summary, the combined Lagrangian resulting from the Higgs mechanism and electroweak symmetry breaking together with the Yukawa terms in short-hand notation is given by

$$\begin{aligned} \mathcal{L}_{\text{Higgs}} = & \underbrace{|\left(\partial_\mu + \frac{1}{2}ig\sigma^a W_\mu^a + \frac{1}{2}ig'Y B_\mu\right)\phi|^2}_{\text{Higgs boson kinetic energy, } W^\pm/Z \text{ boson mass,}} \\ & \text{and } W^\pm/Z \text{ - Higgs boson interaction terms} \\ + & \underbrace{\mu^2|\psi|^2 - \lambda|\psi|^4}_{\text{Higgs boson mass and self-interaction terms}} \\ + & \underbrace{(M_1\bar{\psi}_L\phi\psi_R + M_2\bar{\psi}_L\phi^C\psi_R + h.c.)}_{\text{Yukawa fermion mass and fermion - Higgs interaction terms}} \end{aligned} \quad (2.69)$$

2.1.4 Strong Interaction

The strong interaction between the quarks and gluons of the Standard Model is described by quantum chromodynamics (QCD). The QCD formalism is a local gauge theory based on the symmetry group $SU(3)_C$, where the index C stands for colour charge and is a property of all strongly interacting particles. For quarks the colour charge can take the values “green”, “red”, “blue”, and gluons carry combinations of colour-anticolour. The quarks are assembled in colour triplets of the type

$$\psi = \begin{pmatrix} \psi_r \\ \psi_b \\ \psi_g \end{pmatrix}, \quad (2.70)$$

where the notation

$$\psi_r = \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix}, \quad \psi_b = \begin{pmatrix} 0 \\ 1 \\ 0 \end{pmatrix}, \quad \psi_g = \begin{pmatrix} 0 \\ 0 \\ 1 \end{pmatrix}, \quad (2.71)$$

is adopted. The Lagrangian density for free quarks can then be written just as before

$$\mathcal{L} = \bar{\psi}(i\gamma^\mu \partial_\mu - m)\psi. \quad (2.72)$$

To derive the interaction, gauge invariance is required for local $SU(3)_C$ transformations of the type

$$\psi \rightarrow \psi' = e^{ig_s \alpha^a(x) T^a} \psi, \quad (2.73)$$

where g_s will be identified with the strong coupling constant and the $\alpha^a(x)$ are again arbitrary functions of space and time. The T^a , $a = 1..8$, are the generators of the non-abelian symmetry group $SU(3)$ which follow the commutation relation

$$[T^a, T^b] = if^{abc} T^c, \quad (2.74)$$

where f^{abc} is the so-called structure constant of the symmetry group, given by the tensor with values

$$\begin{aligned} f^{123} &= 1 \\ f^{458} &= f^{678} = \frac{\sqrt{3}}{2} \\ f^{147} &= f^{246} = f^{257} = f^{345} = f^{516} = f^{637} = \frac{1}{2} \end{aligned} \quad (2.75)$$

and antisymmetric permutations thereof, 0 otherwise. A possible representation of the generators $T^a = \lambda_a/2$ are the Gell-Mann matrices

$$\begin{aligned} \lambda_1 &= \begin{pmatrix} 0 & 1 & 0 \\ 1 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, & \lambda_2 &= \begin{pmatrix} 0 & -i & 0 \\ i & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, & \lambda_3 &= \begin{pmatrix} 1 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \\ \lambda_4 &= \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 0 \\ 1 & 0 & 0 \end{pmatrix}, & \lambda_5 &= \begin{pmatrix} 0 & 0 & -i \\ 0 & 0 & 0 \\ i & 0 & 0 \end{pmatrix}, & \lambda_6 &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix}, \\ \lambda_7 &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & -i \\ 0 & i & 0 \end{pmatrix}, & \lambda_8 &= \frac{1}{\sqrt{3}} \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -2 \end{pmatrix}. \end{aligned} \quad (2.76)$$

For the Lagrangian density to stay invariant under $SU(3)_C$ symmetry transformations, the covariant derivative

$$D_\mu = \partial_\mu + ig_s T^a G_\mu^a \quad (2.77)$$

is introduced. It contains eight new vector fields G_μ^a which correspond to the massless gluons. These fields are required to transform as

$$G_\mu^a \rightarrow G_\mu'^a = G_\mu^a - \partial_\mu \alpha^a(x) - g_s f^{abc} \alpha^b(x) G_\mu^c. \quad (2.78)$$

The kinetic energy of the gluons is given by the field strength tensor

$$G_{\mu\nu}^a = \partial_\mu G_\nu^a - \partial_\nu G_\mu^a - g_s f^{abc} G_\mu^b G_\nu^c. \quad (2.79)$$

The complete, gauge invariant Lagrangian density of the strong interactions is thus given by

$$\begin{aligned} \mathcal{L}_{\text{QCD}} &= \bar{\psi}(i\gamma^\mu D_\mu - m)\psi - \frac{1}{4} G^{a\mu\nu} G_{\mu\nu}^a \\ &= \underbrace{\bar{\psi}(i\gamma^\mu \partial_\mu - m)\psi}_{\text{quark kinetic and mass term}} - \underbrace{\frac{1}{2} g_s \bar{\psi}(\gamma^\mu \lambda^a G_\mu^a)\psi}_{\text{quark - gluon field interaction term}} - \underbrace{\frac{1}{4} G^{a\mu\nu} G_{\mu\nu}^a}_{\text{gauge field kinetic and self-interaction term}}. \end{aligned} \quad (2.80)$$

The first term describes the kinetic energy and the mass of the quarks. The second term accounts for the interaction of the quarks with the gluon fields. The third term contains the kinetic energy of the gluon fields and their self interaction with 3 point and 4 point vertices. The self-interacting nature of the gluon fields originates from the fact that the symmetry group $SU(3)_C$ is non-abelian. The self-interaction is also responsible for the confinement of coloured particles, which means that only colourless particles, i.e. mesons consisting of a quark-antiquark pairs, or hadrons consisting of three quarks of different colour may exist freely. Not included in equation 2.80 is an additional CP-violating term inherent to QCD, for which no experimental evidence has been observed to date.

2.1.5 Total Lagrangian of the Standard Model

The combined Lagrangian density of the Standard Model as derived in the previous sections before electroweak symmetry breaking is given by

$$\begin{aligned}
\mathcal{L}_{\text{SM}} = & \underbrace{-\frac{1}{4}B^{a\mu\nu}B_{\mu\nu}^a - \frac{1}{4}W^{a\mu\nu}W_{\mu\nu}^a - \frac{1}{4}G^{a\mu\nu}G_{\mu\nu}^a}_{\text{gauge boson kinetic energy and self-interaction terms}} \\
& + \underbrace{\bar{\psi}_L\gamma^\mu(i\partial_\mu - \frac{1}{2}g\sigma^a W_\mu^a - \frac{1}{2}g'Y B_\mu)\psi_L + \bar{\psi}_R\gamma^\mu(i\partial_\mu - \frac{1}{2}g'Y B_\mu)\psi_R}_{\text{fermion kinetic energy and electroweak interaction terms}} \\
& + \underbrace{|(\partial_\mu + \frac{1}{2}ig\sigma^a W_\mu^a + \frac{1}{2}ig'Y B_\mu)\phi|^2}_{\text{Higgs boson kinetic energy, } W^\pm/Z \text{ boson mass, and } W^\pm/Z - \text{Higgs boson interaction terms}} \\
& + \underbrace{\mu^2|\psi|^2 - \lambda|\psi|^4}_{\text{Higgs boson mass and self-interaction terms}} \\
& + \underbrace{(M_1\bar{\psi}_L\phi\psi_R + M_2\bar{\psi}_L\phi^C\psi_R + h.c.)}_{\text{Yukawa fermion mass and fermion - Higgs interaction terms}} \\
& - \underbrace{\frac{1}{2}g_s\bar{q}\gamma^\mu\lambda^a G_\mu^a q}_{\text{quark - gluon interaction term}}
\end{aligned} \tag{2.81}$$

2.1.6 Experimental Verification of the Standard Model

The Standard Model of particle physics in its simplest form with the assumption of massless neutrinos is defined by 19 parameters: 9 fermion masses, 3 mixing angles and 1 phase from the CKM matrix, 1 strong CP parameter, 3 coupling constants for the electromagnetic, weak, and strong interactions, and 2 boson masses such as e.g. that of the Z and Higgs particles. With the exception of the mass of the Standard Model Higgs boson, 18 out of these 19 parameters have been determined experimentally at some precision and are summarized in reference [2]. With the recent observation of a new particle in the search for the Standard Model Higgs boson the experimental determination of the last missing parameter of the Standard Model seems within close reach. After a review of the theoretical and indirect experimental bounds on the mass of the Higgs boson, the status of the direct searches at the LHC is summarized in this section.

While the mass of the Higgs boson cannot be derived from the known parameters of the Standard Model, some constraints on its upper and lower bounds arise from theoretical arguments. Firstly, the requirement of unitarity of the theory in longitudinal W^+W^- scattering necessitates a Higgs particle with a mass $m_H \lesssim 870$ GeV [28]. This bound can be reduced further with so-called triviality considerations, where the evolution of the quartic Higgs coupling λ as it occurs in eq. 2.57 is examined as a function of the energy scale Q . The behaviour of $\lambda(Q)$ is governed by the renormalization group equations which are described in section 2.2. For low values of $Q^2 \ll v^2$, where v^2 is identified with the electroweak breaking scale, λ converges $\rightarrow 0$. The resulting Higgs potential has no longer the characteristic “Mexican hat” shape shown in Fig. 2.1 and the theory becomes “trivial” since no Higgs self-interactions occur. In the opposite case, $Q^2 \gg v^2$, the quartic coupling eventually becomes infinite and develops a so-called Landau pole at a cut-off energy Λ , where the validity of the Standard Model ends. For a given value of Λ one can then determine an upper limit on the mass of the Higgs boson, e.g. in the case of the Planck scale $\Lambda \sim 10^{16}$ GeV, $m_H \lesssim 200$ GeV is required [29].

Figure 2.2: *Left: Lower and upper bounds on the mass of the Higgs boson M_H from vacuum stability and triviality considerations respectively as a function of the cut-off scale Λ up to which the Standard Model is assumed to be valid. The mass of the top quark and the strong coupling constant are assumed to be at the values given in the figure. The black bands indicate the impact of various uncertainties. Taken from reference [26]. Right: Distribution of $\Delta\chi^2 = \chi^2 - \chi_{\min}^2$ from a global fit to electroweak precision data as a function of different Higgs masses m_H . The preferred value of m_H corresponds to the minimum of the curve at 94_{-24}^{+29} GeV. Values $m_H > 152$ GeV are excluded at 95 % confidence level. The blue band corresponds to the theoretical uncertainty from unknown higher order corrections. Results are shown for two values of $\Delta\alpha_{\text{had}}^{(5)}$, which corresponds to different contributions of the 5 lighter quarks to the QED fine structure constant. The yellow shaded areas show the mass range excluded by the LEP (left) and LHC (right) experiments. Taken from reference [27].*

Lower theoretical bounds on the Higgs mass are derived from vacuum stability considerations, which imply that at low values of λ and m_H contributions from fermions and gauge bosons become significant and can result in negative overall values for the quartic coupling constant. This would imply a Higgs potential without minimum and thus lead to an instability of the vacuum. The lower and upper theoretical bounds on the mass of the Higgs boson from vacuum stability and triviality arguments as a function of the cut off scale Λ are shown in Fig. 2.2 (left).

While theoretical considerations provide a rough window of possible values of the mass of the predicted Higgs particle, experimental measurements are necessary to prove its existence and determine its precise properties. Indirect experimental constraints on m_H arise from global fits to the observables of electroweak precision measurements [27,30], to which radiative corrections from the Standard Model Higgs boson are expected. The current fit results for m_H as a function of $\Delta\chi^2 = \chi^2 - \chi_{\min}^2$ are shown in Fig. 2.2 (right). The preferred value for the mass of the Higgs boson corresponds to the minimum of the curve at a value of 94_{-24}^{+29} GeV. At the upper bound, values $m_H > 152$ GeV can be excluded at 95 % confidence level.

Direct searches for the Higgs boson at the LEP [34], Tevatron [35–37], and LHC [38,39] exper-

Figure 2.3: Top: The observed local p -value is shown for a combination of the considered search channels as a function of possible Higgs boson masses m_H using data from the combined 2011 and 2012 LHC data-taking campaigns at 7 TeV and 8 TeV in proton-proton collision. The dashed line corresponds to the expected local p -value under the hypothesis of a Standard Model Higgs boson. The results of the ATLAS searches are shown on the left and those of the CMS experiment on the right. Bottom left: Measurements of the signal strength parameter μ under the assumption of a Standard Model Higgs boson with $m_H = 125$ GeV for the individual search channels and their combination. Bottom right: Distributions of the log-likelihood ratio generated with pseudo-experiments assuming the spin 0^+ hypothesis and testing the spin 2^- hypothesis in $H \rightarrow ZZ^* \rightarrow 4l$ events. The data corresponds to the solid vertical line. The shaded areas are equivalent to the observed p -values for compatibility with the tested 2^- hypothesis (right shaded area) and the assumed 0^+ hypothesis (left shaded area). Figures taken from references [31–33]

iments have further narrowed the window of possible Higgs masses down to a region between 116 GeV and 127 GeV. In that mass range the ATLAS and CMS experiments have recently reported the observation of a new Higgs-like boson [31, 32, 40, 41]. The statistical significance of the observation corresponds to approximately seven standard deviations as shown on Fig. 2.3

(top). The measured masses are 125.2 ± 0.3 (stat.) ± 0.6 (syst.) GeV and 125.8 ± 0.4 (stat.) ± 0.4 (syst.) GeV for the ATLAS and CMS experiments respectively. According to the measurements performed to date the new particle seems to couple to W and Z bosons as expected, whereas the signal strength in the $H \rightarrow \gamma\gamma$ channel is observed to be somewhat higher as shown in Fig. 2.3 (bottom left). The evidence for couplings to quarks and leptons is weaker at present and will require more data for precise measurements. The measurements of the spin of the new particle are ongoing. Spin-1 can be ruled out since the new particle is observed in di-boson states. As shown in Fig. 2.3 (bottom right) a spin-0 positive parity state (0^+) as expected for the Standard Model Higgs boson seems preferred over 2^- , but the measurements are not fully conclusive yet. A similar tendency is observed with respect to the 0^- and 2^+ hypotheses.

The fact that the mass of the observed Higgs-like particle falls within the narrow window predicted by the fits to electroweak precision data as well as the previously discussed theoretical bounds is a remarkable success of particle physics.

2.2 Supersymmetry

As shown in the previous part of this chapter, the Standard Model of particle physics has been very successful in describing the fundamental particles and their interactions at the currently accessible energy scales in high energy physics. It is expected, however, that more comprehensive theories are necessary to explain the physical phenomena in higher energy regimes. Such theories originate from the Standard Model's inability to clarify important theoretical questions, such as the arbitrariness of gauge couplings, mixing angles and particle masses, the lack of an explanation for gauge symmetry, quantum numbers and generations, the Higgs-mass fine-tuning problem and eventually the incorporation of gravity into a unified theory. In addition significant cosmological observations, such as cold dark matter, dark energy, the observed matter-antimatter asymmetry as well as results from neutrino physics are not addressed within the Standard Model.

The ultimate goal of particle physics is to consistently explain and integrate all these phenomena into a single Theory of Everything (ToE) valid up and beyond the Planck energy scale² $M_P \approx 10^{19} \text{ GeV}$, including quantum gravitational effects. Superstring models [42] are promising candidates on the way to such a ToE, but they are not yet entirely understood and experimental validation is currently beyond technical possibilities.

An essential requirement for most String models is Supersymmetry (SUSY). SUSY is a fundamental symmetry between fermions and bosons introducing a set of new partner particles with opposite spin statistics for each Standard Model particle. The possible implications of SUSY on the electroweak scale make it one of the best-motivated theories beyond the Standard Model. As will be shown in the following, Supersymmetry suggests very elegant solutions to many open questions of the Standard Model.

The Hierarchy Problem. A strong argument for Supersymmetry at the electroweak scale is the so-called Hierarchy Problem [43]. It describes the unnatural discrepancy between the energy scale of the renormalized Higgs-boson mass and that of its bare mass at the lowest order of perturbation theory. This tremendous difference is caused by large quantum corrections to

²The Planck scale is defined as the energy scale at which the effects of gravity become comparable to the other forces and quantum gravity can no longer be ignored.

Figure 2.4: *One-loop quantum correction to the physical Higgs mass from fermion loop (left) and scalar loop (right) diagrams.*

Higgs-boson processes, such as fermion loops as illustrated in Fig. 2.4 (left). These contributions result in quadratically divergent correction terms to the physical Higgs mass:

$$m_H^2 = m_0^2 - \frac{|\lambda_f|^2}{8\pi^2} \Lambda^2 + \mathcal{O}(\ln \frac{\Lambda^2}{m_f^2}) \quad (2.82)$$

Here m_0 is the bare mass of the Higgs-boson at Born-level, λ_f is the Yukawa coupling of the process of interest, m_f the mass of the involved fermion, and Λ a cut-off parameter that can be interpreted as the upper validity limit of the theory and is usually identified with the GUT scale at $\Lambda_{GUT} \approx 10^{16}$ GeV.

In contrast, the mass of the Higgs Boson m_H , as shown in the previous sections, is expected to lie within 116 - 127 GeV. It follows then from equation 2.82 that m_0^2 must be known up to a precision of ~ 24 significant digits to yield the correct value of m_H . From a theoretical point of view it is unlikely that such a fine-tuning of parameters in every order of perturbation theory is realized in nature.

A more elegant solution to this problem is provided by Supersymmetry, where the supersymmetric partner particles, due to their half-spin difference, contribute with opposite sign loop corrections to those of the Standard Model. In this way fermion loop processes cancel with loop diagrams from bosonic SUSY particles as shown in Fig 2.4 (right). Likewise Standard Model boson contributions are absorbed by their respective fermionic superpartners.

For the case of a scalar particle S with mass m_S the corrections can be written as

$$\Delta m_H^2 = \frac{\lambda_S}{16\pi^2} (\Lambda^2 - 2m_S^2 \ln \frac{\Lambda}{m_S} + \dots). \quad (2.83)$$

A comparison between equations 2.82 and 2.83 shows that the quadratic divergences can be absorbed, if there are two scalar contributions and $\lambda_S = |\lambda_f|^2$ holds between the couplings. In fact this relation is an intrinsic property of Supersymmetry as will be shown in the following sections. Now the remaining one loop correction can be approximated as

$$\Delta m_H^2 \approx \mathcal{O}(\frac{\alpha}{\pi})(m_S^2 - m_f^2), \quad (2.84)$$

where α can be identified with a typical coupling constant. Hence, if Supersymmetry is broken,

the masses of the fermions and their supersymmetric partners must lie close together to allow for a natural value of m_H without artificial fine tuning:

$$|m_S^2 - m_f^2| \lesssim 1 \text{ TeV}^2 \quad (2.85)$$

This is one of the strongest motivations to expect Supersymmetry at the electroweak scale.

Grand Unification. Motivated by the evolution of the Standard Model coupling constants with energy, Grand Unified Theories (GUTs) aim to provide a unified description of the electroweak and strong interactions at high energies. The underlying idea is to embed the Standard Model's gauge groups $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ into a universal symmetry group G_{GUT} , representing a single interaction with one coupling constant at a unification scale M_{GUT} .

This requires the three previously defined coupling constants g , g' , and g_s of the SM, if rewritten according to [44]

$$\begin{aligned} \alpha_1 &= (5/3)g'^2/(4\pi) = 5\alpha/(3\cos^2\theta_W), \\ \alpha_2 &= g^2/(4\pi) = \alpha/\sin^2\theta_W, \\ \alpha_3 &= g_s^2/(4\pi), \end{aligned} \quad (2.86)$$

to intersect at M_{GUT} .

The energy dependence of the parameters α_i is provided by the renormalization formalism. Using a specific renormalization scheme³, the contributions of vacuum polarization processes to the boson propagators are taken into account. The corresponding renormalization group equations at one-loop level are given by [47]

$$\frac{d\alpha_i}{dt} = \frac{b_i}{2\pi}\alpha_i^2, \quad i = 1, 2, 3, \quad (2.87)$$

with $t = \ln(Q^2/\mu_R^2)$, where Q^2 stands for the energy of the interaction and μ_R^2 is the renormalization scale. The analytical solution of equation (2.87),

$$\frac{1}{\alpha_i(Q^2)} = \frac{1}{\alpha_i(\mu_R^2)} - \frac{b_i}{2\pi} \ln\left(\frac{Q^2}{\mu_R^2}\right), \quad (2.88)$$

describes the so-called 'running' of the coupling constants as a function of energy. The coefficients b_i carry intrinsic information about the particle content of the underlying model. In the case of the Standard Model they have been determined to [48]

$$b_i = \begin{pmatrix} b_1 \\ b_2 \\ b_3 \end{pmatrix} = \begin{pmatrix} 0 \\ -22/3 \\ -11 \end{pmatrix} + N_{Fam} \begin{pmatrix} 4/3 \\ 4/3 \\ 4/3 \end{pmatrix} + N_{Higgs} \begin{pmatrix} 1/10 \\ 1/6 \\ 0 \end{pmatrix}, \quad (2.89)$$

where $N_{Fam} = 3$ stands for the number of generations and $N_{Higgs} = 1$ is the number of Higgs doublets.

³Here the so-called modified minimal subtraction scheme $\overline{MS}$ is used [45, 46].

Figure 2.5: *Evolution of the inverse coupling constants with energy. Within the Standard Model (left) a unification is not natural, whereas for minimal supersymmetric models (right) the coupling constants intersect at $M_{GUT} \approx 10^{16}$ GeV [44]. The thickness of the lines indicates the error on the coupling constants. The plots are taken from [49].*

Using equations (2.88), (2.89) and the experimentally measured values for the couplings, one can extrapolate to high energies to examine a possible unification [44]. The result of this extrapolation is shown in Fig. 2.5 (left), which clearly indicates that a unification in a single point is not natural. In fact it is ruled out by more than 7 standard deviations [44] and so is a minimal GUT based on the Standard Model.

To maintain the idea of a Grand Unified Theory, new physics must exist between the electroweak and Planck scale to alter the behaviour of the α_i . This is where Supersymmetry enters the picture. Assuming a minimal supersymmetric model as described in section 2.2.3, the extended particle content changes the coefficients b_i to [48]

$$b_i = \begin{pmatrix} b_1 \\ b_2 \\ b_3 \end{pmatrix} = \begin{pmatrix} 0 \\ -6 \\ -9 \end{pmatrix} + N_{Fam} \begin{pmatrix} 2 \\ 2 \\ 2 \end{pmatrix} + N_{Higgs} \begin{pmatrix} 3/10 \\ 1/2 \\ 0 \end{pmatrix}, \quad (2.90)$$

where $N_{Fam} = 3$ remains and an additional Higgs doublet is introduced $N_{Higgs} = 2$. It turns out, with these parameters a unification becomes possible at $\sim 10^{16}$ GeV as illustrated in Fig. 2.5 (right). Furthermore the most perfect intersection of the coupling constants can be obtained if the masses of the SUSY particles are of the order of 1 TeV [44]. This provides yet another very strong motivation for low energy Supersymmetry.

Dark Matter. A third argument in favour of Supersymmetry is its possible explanation for dark matter. Dark matter is a hypothetical form of matter that cannot be observed directly, but whose existence can be inferred from gravitational effects on visible matter, such as the rotation of galaxies or the structure formation in the universe. Latest cosmological results, e.g. from WMAP-data [50], suggest that the cold dark matter content makes up about 23% of the

energy density of the universe.

Candidate particles for dark matter include the so-called weakly-interacting massive particles (WIMPs). These hypothetical neutral particles only participate in the gravitational and weak interactions and thus are extremely difficult to detect. Neutrinos are the only WIMP-like particles within the Standard Model. They are, however, not massive and abundant enough to provide an explanation for dark matter.

In R-Parity conserving SUSY models (see section 2.2.3) the lightest supersymmetric particle (LSP) is stable and can be neutral and only weakly interacting. It therefore exhibits all features of a WIMP and constitutes a possible dark matter candidate.

Gravity. As a last important point the connection between Supersymmetry and Gravity is briefly mentioned. As will be further explained in section 2.2.5, by making Supersymmetry a local symmetry the principles of both theories can be unified into a single concept, called Supergravity.

Supergravity is non-renormalizable and therefore not a candidate for a Theory of Everything, but it can be understood as an effective description of physical phenomena including gravity at energies below the Planck scale [51].

In conclusion, a variety of reasons point to Supersymmetry as a possible theory for physics beyond the Standard Model. In the following sections of this chapter the mathematical formalism of Supersymmetry is introduced step-by-step.

2.2.1 Supersymmetry Algebra

In this section the group-theoretical foundations of Supersymmetry are reviewed. As a starting point the so-called Poincaré group is considered, one of the most fundamental symmetry groups in physics. It contains the full symmetry of special relativity, including the four translations in Minkowski space, plus the three rotations and boosts of the Lorentz group. The known elementary particles are irreducible representations of the Poincaré group.

To obtain a unified theory of all interactions, it is desirable to combine the internal gauge symmetries of the Standard Model, represented by the Lie groups $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$, with the structure of space-time provided by the Poincaré group. Unfortunately, the existence of such a group is ruled out for any but the trivial case, as was shown by Coleman and Mandula [52] in 1967. However, after Wess and Zumino found the first supersymmetric model in 1974 [53], this theorem had to be significantly generalized. In fact it was shown by Haag, Lopuszanski, and Sohnius [54] a year later that under weaker assumptions a non-trivial extension to the Poincaré symmetry is possible, namely *Supersymmetry*.

Thus Supersymmetry is a space-time symmetry. Its generators Q transform bosonic states into fermionic states and vice-versa:

$$Q|Boson\rangle = |Fermion\rangle, \quad Q|Fermion\rangle = |Boson\rangle \quad (2.91)$$

It follows that Q and its hermitian conjugate $Q^\dagger$ must have fermionic character and carry spin-1/2. Within the easiest supersymmetric extension of the Poincaré group, the above mentioned Wess-Zumino model, the generators are required to satisfy the following algebra,

$$\{Q_\alpha, Q_{\dot{\alpha}}^\dagger\} = 2\sigma_{\alpha\dot{\alpha}}^\mu P_\mu \quad (2.92)$$

$$\{Q_\alpha, Q_\beta\} = \{Q_{\dot{\alpha}}^\dagger, Q_{\dot{\beta}}^\dagger\} = 0 \quad (2.93)$$

$$[Q_\alpha, P^\mu] = [Q_{\dot{\alpha}}^\dagger, P^\mu] = 0, \quad (2.94)$$

where P^μ denotes the four-momentum in space-time, σ^μ are the Pauli matrices and $\alpha, \dot{\alpha}, \beta, \dot{\beta}$ are the indices of two-component Weyl spinors.

Direct implications of this supersymmetric algebra are the following:

- The bosonic or fermionic states and their respective superpartners with opposite spin statistics are ordered in supersymmetric multiplets, so-called *supermultiplets*. The supermultiplets are the irreducible representations of the supersymmetric algebra.
- Superpartners must have equal mass, since $-P^2$ commutes with both Q and $Q^\dagger$ as well as all generators of the Poincaré group.
- Superpartners must have equal gauge quantum numbers, such as charge, isospin and color, since Q and $Q^\dagger$ commute with all generators of the internal gauge symmetry groups.
- A concatenation of two supersymmetric transformations leads to a translation in space-time, since the square of the generator Q is equal to P_μ . For local supersymmetric transformations this implies a connection between Supersymmetry and General Relativity as will be shown later on.
- The bosonic and fermionic degrees of freedom of each supermultiplet are related as $n_B = n_F$.

The latter statement needs a short explanation: As stated earlier the fermionic generators of Supersymmetry Q and $Q^\dagger$ map the bosonic subspace B onto a fermionic subspace F and vice-versa. It is also known that for a linear mapping $f : X \rightarrow Y$ the relation $\dim(Y) \leq \dim(X)$ holds. Thus, for a two-fold supersymmetric transformation $B \rightarrow F \rightarrow B$, the equation $\dim(F) = \dim(B)$ must apply, since the concatenation of two SUSY transformations maps the bosonic subspace onto itself. The same argument holds for the fermionic subspace and thus the bosonic and fermionic degrees of freedom in each supermultiplet must be equal.

Following this rule the possible constellations of supermultiplets are examined. The easiest example is the so-called *chiral supermultiplet*. Each chiral supermultiplet consists of one fermion with two spin degrees of freedom $n_F = 2$ and two real scalar fields with $n_B = 1$ each. The two real components are equivalent to one complex scalar field. Per naming convention, the supersymmetric scalar particle states receive an "s"-prefix to their name ("s-fermion"), to distinguish them from the original particle.

Next are the gauge supermultiplets. They contain massless spin-1 vector bosons with two helicity states and again a spin-1/2 fermion with $n_F = 2$ as superpartner. Here the superparticles are indicated by an "ino"-suffix to the name of the corresponding gauge boson ("gaugino").

Depending on the underlying supersymmetric model there are also other possible constellations. An explicit discussion of the supermultiplets and their particle content will be given in section 2.2.3 for the case of the Minimal Supersymmetric Standard Model.

2.2.2 Supersymmetric Lagrangians

In this section the Lagrangian formalism of Supersymmetry is derived. The goal is to obtain a general formulation of the supersymmetric field theory with its particle masses and interactions. Following closely the approach in [43], the Lagrangians of the chiral and gauge supermultiplets and their respective field interactions are introduced step-by-step.

The Chiral Supermultiplet

To begin with, the simplest possible SUSY model, containing a single left-handed two component Weyl fermion ψ and a complex scalar field ϕ as its superpartner is considered. This is equivalent to a theory with only one chiral supermultiplet, also referred to as the massless non-interacting Wess-Zumino model [53]. The Lagrangian density of this model is of the form

$$\mathcal{L}_{\text{chiral,free}} = \mathcal{L}_{\text{scalar}} + \mathcal{L}_{\text{fermion}} = -\partial^\mu \phi^* \partial_\mu \phi - i\psi^\dagger \bar{\sigma}^\mu \partial_\mu \psi, \quad (2.95)$$

and contains two kinetic terms for the scalar and fermionic states.

Now minimal supersymmetric transformations of the scalar fields of the type

$$\delta\phi = \epsilon\psi \quad \text{and} \quad \delta\phi^* = \epsilon^\dagger\psi^\dagger \quad (2.96)$$

are introduced, where ϵ is an infinitesimal fermion-like parameter. The scalar part of the Lagrangian transforms under these equations as

$$\delta\mathcal{L}_{\text{scalar}} = -\epsilon\partial^\mu\psi\partial_\mu\phi^* - \epsilon^\dagger\partial^\mu\psi^\dagger\partial_\mu\phi. \quad (2.97)$$

According to Hamilton's principle the action (eq. 2.1) must stay invariant under any symmetry transformations. Therefore the Lagrangians before and after the transformation must be equivalent up to a total derivative $\mathcal{L}' = \mathcal{L} + \partial_\mu\Lambda^\mu$. This implies the following transformations of the fermion fields:

$$\delta\psi_\alpha = i(\sigma^\mu\epsilon^\dagger)_\alpha\partial_\mu\phi \quad \text{and} \quad \delta\psi^\dagger_{\dot{\alpha}} = -i(\epsilon\sigma^\mu)_{\dot{\alpha}}\partial_\mu\phi^* \quad (2.98)$$

Using Pauli matrix identities and commutation relations for partial derivatives one obtains

$$\delta\mathcal{L}_{\text{fermion}} = \epsilon\partial^\mu\psi\partial_\mu\phi^* + \epsilon^\dagger\partial^\mu\psi^\dagger\partial_\mu\phi - \partial_\mu(\epsilon\sigma^\nu\bar{\sigma}^\mu\psi\partial_\nu\phi^* + \epsilon\psi\partial^\mu\phi^* + \epsilon^\dagger\psi^\dagger\partial^\mu\phi). \quad (2.99)$$

Comparing 2.97 and 2.99 one sees that the two contributions cancel up to a total derivative and thus leave the action invariant as required.

However, for the SUSY algebra to be valid off-shell, this formalism needs to be generalized by the introduction of an auxiliary field F . F has no kinematic term and its Lagrangian density is given by

$$\mathcal{L}_{\text{auxiliary}} = F^*F. \quad (2.100)$$

The auxiliary fields transform as

$$\delta F = i\epsilon^\dagger\bar{\sigma}^\mu\partial_\mu\psi \quad \text{and} \quad \delta F^* = -i\partial_\mu\psi^\dagger\bar{\sigma}^\mu\epsilon, \quad (2.101)$$

and to keep the action invariant equations 2.98 have to be modified with additional F -terms

$$\delta\psi_\alpha = i(\sigma^\mu\epsilon^\dagger)_\alpha\partial_\mu\phi + \epsilon_\alpha F \quad \text{and} \quad \delta\psi^\dagger_{\dot{\alpha}} = -i(\epsilon\sigma^\mu)_{\dot{\alpha}}\partial_\mu\phi^* + \epsilon^\dagger_{\dot{\alpha}}F^*. \quad (2.102)$$

The need for the auxiliary field F becomes apparent, if one compares the scalar and fermionic degrees of freedom on- and off-shell. In the on-shell case, the complex scalar field ϕ has two real components corresponding to the two helicity states of the fermion field ψ . Going off-shell, however, ψ becomes a two-dimensional complex object with four real degrees of freedom. To balance out this inequality the complex field F with two additional degrees of freedom must be introduced.

The final Lagrangian of the free chiral supermultiplet can now be written as

$$\boxed{\mathcal{L}_{\text{chiral,free}} = -\partial^\mu\phi^{*i}\partial_\mu\phi_i - i\psi^\dagger\bar{\sigma}^\mu\partial_\mu\psi_i + F^{*i}F_i}, \quad (2.103)$$

where i is the index over all flavour and gauge degrees of freedom.

Interactions. The interactions between the scalar and fermion fields inside the supermultiplets are now considered. It can be shown [43] that the most general renormalizable form of the interaction Lagrangian is given by

$$\mathcal{L}_{\text{chiral,int}} = \left(-\frac{1}{2}W^{ij}\psi_i\psi_j + W^iF_i\right) + c.c. , \quad (2.104)$$

where W^{ij} and W^i are functions of the scalar fields participating in the interaction. The invariance of $\mathcal{L}_{\text{chiral,int}}$ under SUSY transformations implies the following form of the W^{ij} and W^i [43]

$$W^{ij} = \frac{\partial^2}{\partial\phi_i\partial\phi_j}W = \frac{1}{2}M^{ij} + \frac{1}{6}y^{ijk}\phi_k, \quad (2.105)$$

$$W^i = \frac{\partial W}{\partial\phi_i} = \frac{1}{2}M^{ij}\phi_j + \frac{1}{6}y^{ijk}\phi_j\phi_k, \quad (2.106)$$

where W is the so-called superpotential

$$W = \frac{1}{2}M^{ij}\phi_i\phi_j + \frac{1}{6}y^{ijk}\phi_i\phi_j\phi_k. \quad (2.107)$$

Here M^{ij} is the symmetric fermion mass matrix and y^{ijk} the Yukawa coupling of two fermion fields with one scalar. It can be seen that the superpotential contains only bilinear and trilinear scalar coupling terms and no fermionic contributions.

The auxiliary field terms of Lagrangians 2.103 and 2.104 are $F_iF^{i*} + W^iF_i + W_i^*F^{i*}$. They can be reformulated in terms of the superpotential by applying the Euler-Lagrange equations of motion for F^i leading to

$$F_i = -W_i^* \quad \text{and} \quad F^{i*} = -W^i. \quad (2.108)$$

With this the auxiliary-free representation of the total Lagrangian of the interacting chiral supermultiplet is obtained:

$$\mathcal{L}_{chiral} = -\partial^\mu \phi^{i*} \partial_\mu \phi_i - V(\phi, \phi^*) - i\psi^{i\dagger} \bar{\sigma}^\mu \partial_\mu \psi_i - \frac{1}{2} M^{ij} \psi_i \psi_j - \frac{1}{2} M_{ij}^* \psi^{i\dagger} \psi^{j\dagger} - \frac{1}{2} y^{ijk} \phi_i \psi_j \psi_k - \frac{1}{2} y_{ijk}^* \phi^{i*} \psi^{j\dagger} \psi^{k\dagger} \quad (2.109)$$

Here V is referred to as the scalar potential

$$V(\phi, \phi^*) = W^k W_k^* = F^{*k} F_k = M_{ik}^* M^{kj} \phi^{i*} \phi_j + \frac{1}{2} M^{in} y_{jkn}^* \phi_i \phi^{j*} \phi^{k*} + \frac{1}{2} M_{in}^* y^{jkn} \phi^{i*} \phi_j \phi_k + \frac{1}{4} y^{ijn} y_{kln}^* \phi_i \phi_j \phi^{k*} \phi^{l*}. \quad (2.110)$$

It can be seen that the scalar potential contains cubic and quartic scalar coupling terms as well as a mass term with the same mass matrix as the fermionic part in equation 2.109. This leads to the expected mass-degenerate partner states inside each supermultiplet.

Equation 2.109 also demonstrates that the coupling strength of a scalar particle with two fermion fields is of the order y^{ijk} , whereas equation 2.110 implies $(y^{ijk})^2$ for a quartic scalar process. This is equivalent to the relation $\lambda_S = \lambda_f^2$ that was postulated in the introduction of section 2.2 for the solution of the hierarchy problem.

The Gauge Supermultiplet

The gauge supermultiplets consist of massless gauge bosons A_μ^a and their gaugino superpartners λ_a

$$\begin{pmatrix} A_\mu^a \\ \lambda_a \end{pmatrix}, \quad (2.111)$$

where a is the index over the group representations, e.g. $a = 1, \dots, 8$ for $SU(3)_C$. The corresponding supersymmetric transformations are given by

$$\delta_{gauge} A_\mu^a = \partial_\mu \Lambda^a + g f^{abc} A_\mu^b \Lambda^c \quad (2.112)$$

$$\delta_{gauge} \lambda^a = g f^{abc} \lambda^b \Lambda^c. \quad (2.113)$$

Here Λ is an infinitesimal gauge transformation parameter, g represents the coupling strength of the interaction and f^{abc} the structure constant in the case of a non-abelian theory. The total Lagrangian density of the gauge supermultiplets must again leave the action invariant under equations 2.112 and 2.113. It is given by:

$$\mathcal{L}_{gauge} = -\frac{1}{4} F^{a\mu\nu} F_{\mu\nu}^a - i\lambda^{a\dagger} \bar{\sigma}^\mu D_\mu \lambda^a + \frac{1}{2} D^a D^a \quad (2.114)$$

The first term describes the kinetic energy with the gauge field tensor

$$F_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a + g f^{abc} A_\mu^b A_\nu^c, \quad (2.115)$$

where the A_μ are the respective gauge field components of the interaction. The second term in equation 2.114 expresses the kinetic energy of the gaugino fields and includes a covariant derivative

$$D_\mu \lambda^a = \partial_\mu \lambda^a + g f^{abc} A_\mu^b \lambda^c, \quad (2.116)$$

which contains the interactions between the gauge and gaugino fields. With the third term in equation 2.114 again an auxiliary field D is introduced to account for the inequality of degrees of freedom within the supermultiplets on- and off-shell. In this case only one additional degree of freedom is necessary and thus D has one real component. As in the case of the chiral Lagrangian this field vanishes when going on-shell.

The Combined Supersymmetric Lagrangian

Now the results for the Lagrangian of the chiral and gauge supermultiplets can be combined. First, the derivatives of the scalar and fermion fields need to be replaced by the respective covariant derivative to preserve gauge invariance:

$$\partial_\mu \psi \rightarrow D_\mu \psi = \partial_\mu \psi + ig A_\mu^a T^a \psi \quad (2.117)$$

$$\partial_\mu \phi \rightarrow D_\mu \phi = \partial_\mu \phi + ig A_\mu^a T^a \phi \quad (2.118)$$

Here T^a stands for the generators of the gauge groups, and one can see that equations 2.117 and 2.118 yield the couplings between the gauge bosons and the scalar and fermionic fields of the chiral supermultiplet.

There are also possible couplings between the gaugino fields λ^a and the D auxiliary fields. It can be shown [43] that possible renormalizable couplings of this sort contribute with terms:

$$-\sqrt{2}g(\phi^* T^a \psi) \lambda^a - \sqrt{2}g \lambda^{\dagger a} (\phi^\dagger T^a \phi) + g(\psi^* T^a \phi) D^a \quad (2.119)$$

The D -term in equation 2.114 and the last term in 2.119 combine to the equation of motion

$$D^a = -g(\phi^* T^a \phi). \quad (2.120)$$

Since 2.120 contains only scalar fields it is then usually written with the scalar potential

$$V(\phi, \phi^*) = W_i^* W^i + \frac{1}{2} g^2 (\phi^* T \phi)^2. \quad (2.121)$$

The results of this chapter can now be summarized in the total Lagrangian density for Supersymmetry:

$$\begin{aligned}
\mathcal{L}_{SUSY} = & \underbrace{-D^\mu \phi^* D_\mu \phi}_{\text{scalars}} \underbrace{-i\psi^\dagger \bar{\sigma}^\mu D_\mu \psi}_{\text{fermions}} \underbrace{-\frac{1}{2}(W^{ij}\psi_i\psi_j + W^{ij*}\psi^\dagger_i\psi^\dagger_j)}_{\text{fermion mass term and Yukawa coupling}} \\
& \underbrace{-W^i W_i^* - \frac{1}{2}g^2(\phi^* T \phi)^2}_{\text{scalar potential}} \underbrace{-\frac{1}{4}F_{\mu\nu}^a F^{a\mu\nu}}_{\text{gauge bosons}} \underbrace{-i\lambda^{\dagger a} \bar{\sigma}^\mu D_\mu \lambda^a}_{\text{gauginos}} \\
& \underbrace{-\sqrt{2}g(\phi^* T^a \psi)\lambda^a - \sqrt{2}g\lambda^{\dagger a}(\psi^\dagger T^a \phi)}_{\text{additional couplings}} \underbrace{+\mathcal{L}_{\text{soft}}}_{\text{soft breaking term}}
\end{aligned} \tag{2.122}$$

The first two terms are the kinetic terms of the scalar and fermionic fields in the chiral supermultiplet. They are followed by the fermion mass terms and the Yukawa coupling of scalar and fermionic fields, written in terms of the superpotential 2.107. The mass terms of the scalar fields and their associated interactions are contained in the scalar potential. The next three terms account for the kinetic energy and interactions of the gauge bosons and their gaugino superpartners plus additional couplings as discussed above. For completeness the SUSY breaking term $\mathcal{L}_{\text{soft}}$ is also included and will be discussed in section 2.2.4.

2.2.3 The Minimal Supersymmetric Standard Model

In the previous section of this chapter the general theoretical features of Supersymmetry have been discussed. Now the Minimal Supersymmetric Standard Model (MSSM) is discussed specifically. It represents the simplest possible supersymmetric extension of the Standard Model and has long been the main focus of experimental searches. Implications of a possible Higgs boson in the mass range of 126 GeV and recent experimental bounds on the MSSM will be discussed in section 2.2.7.

Particle content of the MSSM

The MSSM contains the minimal number of couplings and fields. The field content is described in terms of:

- Chiral supermultiplets with SM leptons and quarks and their associated scalar superpartners ("squarks" and "sleptons").
- Gauge supermultiplets with SM gauge bosons and associated superpartners ("gauginos").

A summary of all MSSM multiplets can be found in Tables 2.2 and 2.3. For the chiral supermultiplets, each SM fermion has two helicity states, which transform differently under gauge symmetry. Therefore they must both have their own scalar superpartner. This is illustrated in Table 2.2, e.g. for the left- and right-handed electrons e_L and e_R and their corresponding scalar particles $\tilde{e}_L$ and $\tilde{e}_R$.⁴ The latter two scalar fields are completely independent and the index refers only to the handedness of the associated SM particles. All fermions in the chiral

⁴The tilde symbol is used to denote supersymmetric partners of SM particles.

supermultiplets are defined in terms of left-handed Weyl-spinors. Conjugations are therefore applied to the right-handed fields of Table 2.2.

In the supersymmetric Higgs-sector, the scalar Higgs fields are accommodated in the chiral supermultiplets along with their "Higgsino" superpartners (see Table 2.2). However, within the MSSM one Higgs-doublet is not sufficient. In particular, gauge anomalies from triangle diagrams, as they are known from the SM, do not cancel within the MSSM unless a second doublet is introduced. Two Higgs-doublets are also required to give mass to all matter fermions by means of electroweak symmetry breaking. The doublets are of the form

$$\begin{pmatrix} H_u^+ \\ H_u^0 \end{pmatrix} \quad \text{and} \quad \begin{pmatrix} H_d^0 \\ H_d^- \end{pmatrix}, \quad (2.123)$$

with weak Isospin $Y = \pm 1/2$ to induce the necessary Yukawa couplings to all up- and down-like quarks. As can be seen from equation 2.123 there are 4 complex or 8 real degrees of freedom in the Higgs doublets. As in the Standard Model three phases are absorbed by Goldstone bosons leaving 5 physical Higgs eigenstates:

- h , a light neutral scalar Higgs particle.
- H , a heavy neutral scalar Higgs Particle.
- A , a neutral CP-odd pseudoscalar Higgs particle.
- $H^\pm$, two charged scalar Higgs particles.

The gauge supermultiplets are shown in Table 2.3. They contain the mediators of the SM interactions and their spin-1/2 superpartners. These are 8 gluons and gluinos in the case of $SU(3)_C$ for QCD and $W^\pm, W^0, B^0$ with superpartners for $SU(2)_L \otimes U(1)_Y$. After electroweak symmetry breaking the latter four mix to the mass eigenstates Z^0 and γ and the respective gaugino combinations "Zino" $\tilde{Z}^0$ and "Photino" $\tilde{\gamma}$.

Electroweak gauginos also form new mass eigenstates with the Higgsinos of same charge. In the neutral sector $\tilde{W}^0, \tilde{B}^0, \tilde{H}_u^0, \tilde{H}_d^0$ mix to four "Neutralinos" $\tilde{\chi}_{1,2,3,4}^0$, whereas the charged $\tilde{W}^\pm, \tilde{H}_u^\pm$ and $\tilde{H}_d^\pm$ form the "Charginos" $\tilde{\chi}_{1,2}^\pm$. The mixing is possible because the participating states have identical quantum numbers as can be inferred from Table 2.3.

It should be noted here that if Supergravity (see section 2.2.5) is to be included into the theory, one must also introduce an additional type of supermultiplet containing the spin-2 graviton and its spin-3/2 superpartner ("gravitino").

R-parity

An important quantum number in the MSSM is R-parity. Whereas in the Standard Model baryon- and lepton-numbers are automatically conserved, the MSSM theoretically allows interaction terms that violate this symmetry. To avoid this undesired effect, the conservation of R-parity

$$R = (-1)^{3(B-L)+2S} \quad (2.124)$$

Name		Scalar ϕ (S=0)	Fermion ψ (S=1/2)
Squarks, Quarks (3 Generations)	Q_1	$(\tilde{u}_L, \tilde{d}_L)$	(u_L, d_L)
	$\bar{U}_1$	$\tilde{u}_R^*$	$u_R^\dagger$
	$\bar{D}_1$	$\tilde{d}_R^*$	$d_R^\dagger$
Sleptons, Leptons (3 Generations)	L_1	$(\tilde{\nu}_e, \tilde{e}_L)$	(ν_e, e_L)
	$\bar{E}_1$	$\tilde{e}_R^*$	$e_R^\dagger$
Higgs, Higgsino	H_u	(H_u^+, H_u^0)	$(\tilde{H}_u^+, \tilde{H}_u^0)$
	H_d	(H_d^0, H_d^-)	$(\tilde{H}_d^0, \tilde{H}_d^-)$

Table 2.2: *Content of the chiral supermultiplets in the Minimal Supersymmetric Standard Model. The symbols Q_i and L_i stand for the supermultiplets containing $SU(2)_L$ doublets, while $\bar{U}_i$, $\bar{D}_i$, and $\bar{E}_i$ contain the corresponding conjugate right-handed singlet states. Table adapted from [43].*

is imposed. B and L are lepton and baryon numbers and S the spin of the participating particles in the process. All SM fields carry R-parity $R = +1$, all superfields $R = -1$. In addition R-parity is a multiplicative quantum number. It follows that in this case supersymmetric particles can only be produced in pairs by SM particles due to

$$R_{total} = R_1^{SM} \cdot R_2^{SM} = 1^2 = R_1^{SUSY} R_2^{SUSY} = (-1)^2 = 1. \quad (2.125)$$

Another consequence of R-parity conservation is the stability of the lightest supersymmetric particle. Due to $R = -1$ it cannot decay into SM matter and due to its mass not into any other supersymmetric particle. It is therefore considered a candidate for dark matter.

In some supersymmetric models, which are not the subject of this thesis, the violation of R-parity is allowed to the extent that it is compatible with the observed proton lifetime.

Interactions of the MSSM

In the MSSM the generic superpotential of equation 2.107 is replaced by [43]

$$W_{MSSM} = y_u \bar{U} Q H_U - y_d \bar{D} Q H_d - y_e \bar{E} L H_d + \mu H_u H_d, \quad (2.126)$$

where Q , L , $\bar{U}$, $\bar{D}$, $\bar{E}$, H_u , and H_d stand for the superfields of the respective multiplets shown in table 2.2. The indices for each generation of quarks and leptons are suppressed in this vector notation. The 3×3 matrices y_u , y_d and y_e are the corresponding Yukawa couplings and determine the CKM mixing angles and masses after electroweak symmetry breaking. The last term is the supersymmetric Higgs term with mass parameter μ .

The supersymmetric Yukawa interactions can be derived from eq. 2.126. They describe cubic and quartic couplings of fermions, sfermions, Higgs, and Higgsino fields. The coupling of SM gauge bosons to supersymmetric particles is governed by the kinetic terms of the SUSY Lagrangian. The MSSM gauge interactions can be obtained by interchanging any two of the three participating particles in a SM gauge interaction by their respective superpartners. In

Name	Boson A^μ (S=1)	Fermion λ (S=1/2)
Gluon, Gluino	g	$\tilde{g}$
W-Bosons, Winos	$W^\pm, W^0$	$\tilde{W}^\pm, \tilde{W}^0$
B-Boson, Bino	B^0	$\tilde{B}^0$

Table 2.3: *Content of the gauge supermultiplets in the Minimal Supersymmetric Standard Model. Table adapted from [43].*

this way one can also obtain gaugino-sfermion-fermion interactions, which are possible through the additional renormalizable couplings, corresponding to the first two terms in equation 2.119.

2.2.4 Breaking of Supersymmetry

If Supersymmetry exists it must be broken, since no mass-degenerate superpartners of the SM particles have been found. The breaking mechanism should preserve the renormalizability of the theory as well as the cancellation of quadratic divergences to maintain the hierarchy of the energy scales as discussed in the introduction. A symmetry breaking with these basic SUSY properties can be introduced into the theory by adding a so-called "soft-breaking" term in the Lagrangian density. The general form of this term is [43]

$$\mathcal{L}_{\text{soft}} = -\left(\frac{1}{2}M_a\lambda^a\lambda^a + \frac{1}{6}a^{ijk}\phi_i\phi_j\phi_k + \frac{1}{2}b^{ij}\phi_i\phi_j + c.c.\right) - (m^2)_j^i\phi^{j*}\phi_i, \quad (2.127)$$

with squared scalar mass terms $(m^2)_j^i$ and b^{ij} , cubic scalar couplings a^{ijk} , and gaugino mass terms M_a for each gauge group. It can be seen that 2.127 contains only scalar and gaugino terms and thus breaks the symmetry by giving masses to the associated particles. A phenomenological explanation of this breaking mechanism will follow in the next section. In the case of the MSSM the soft breaking term specializes to [43]

$$\begin{aligned} \mathcal{L}_{\text{soft}}^{MSSM} = & -\frac{1}{2}(M_3\tilde{g}\tilde{g} + M_2\tilde{W}\tilde{W} + M_1\tilde{B}\tilde{B} + c.c.) \\ & -(\bar{U}a_UQH_u - \bar{D}a_DQH_d - \bar{E}a_ELH_d + c.c.) \\ & -Q^\dagger m_Q^2 Q - L^\dagger m_L^2 L - \bar{U}m_U^2 \bar{U}^\dagger - \bar{D}m_D^2 \bar{D}^\dagger - \bar{E}m_E^2 \bar{E}^\dagger \\ & -m_{H_u}^2 H_u^* H_u - m_{H_d}^2 H_d^* H_d - (bH_u H_d + c.c.). \end{aligned} \quad (2.128)$$

The first line represents the mass terms of winos, gluinos and binos, and the third line those of the squarks and sleptons with hermitian 3×3 mass matrices. Here the tildes on the scalar superfields $Q, \bar{U}, \bar{D}, L, \bar{E}$ are again suppressed for readability. The second line contains the cubic scalar couplings of equation 2.127 with the matrices a_U, a_D, a_E , where again all three generations contribute. The last line corresponds to the soft breaking contributions from the squared Higgs-mass terms $m_{H_u}^2$ and $m_{H_d}^2$ plus one b^{ij} -type term.

Equation 2.128 demonstrates the complexity of the spontaneously broken MSSM. In total 105 new parameters are introduced: 21 masses, 36 mixing angles, 40 CP-violating phases in the

squark and slepton sector and 5 real and 3 CP-violating parameters in the Higgs-sector. However, not all of these parameters are independent. In particular flavour and CP-conserving relations reduce the number of degrees of freedom significantly.

The most common framework for SUSY breaking models is based on a so-called "hidden sector", in which the symmetry is spontaneously broken [55] [56]. The introduction of this hidden sector is necessary, since none of the MSSM fields, which are referred to as the "visible sector", can have a non-zero vacuum expectation value in order to not violate gauge invariance⁵. Therefore the underlying idea of the hidden sector is, that the spontaneous SUSY breaking is communicated down to the observable MSSM sector via hypothetical flavour-blind messenger fields. This mediation mechanism, however, is highly dependent on the model framework assumed.

Some of the most common scenarios are:

- Gauge Mediated Supersymmetry Breaking (GMSB)
- Anomaly Mediated Supersymmetry Breaking (AMSB)
- Gravity Mediated Supersymmetry Breaking (MSUGRA)

In GMSB the SUSY breaking is invoked by the minimal gauge group interactions $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$, where the associated gauge bosons couple to the messenger fields to make the breaking observable. AMSB is a special case of gravity-mediated SUSY breaking, where there are no direct tree-level couplings between the sectors and the masses of the sparticles are generated with higher order loop corrections. MSUGRA is one of the most commonly assumed scenarios in experimental searches. It will be discussed in more detail in the following section.

2.2.5 Minimal Supergravity

The principles of Supergravity were already briefly mentioned in the introduction of section 2.2. The underlying idea is that supersymmetry, when made a local symmetry, yields both, an effective field theory for energies below the Planck scale and an elegant mechanism for SUSY breaking. In this way Supergravity, as any local gauge theory, necessitate the introduction of new gauge fields, the spin-2 graviton and its spin-3/2 gravitino superpartner. In an unbroken theory the masses of both particles are zero. The couplings, however, still scale with the dimensional Newton's constant and are thus proportional to $\sim 1/M_{\text{Planck}}$. Therefore the associated terms in the Lagrangian are non-renormalizable and so is the concept of Supergravity, which hence does not represent a full theory of quantum gravity.

The breaking mechanism of Supergravity, also referred to as super-Higgs-mechanism, takes place in two steps: First the spontaneous breaking of *global* Supersymmetry yields a massless Weyl fermion called Goldstino. This Goldstino has two degrees of freedom which are subsequently absorbed through the spontaneous breaking of *local* Supersymmetry to give mass to the gravitino. The graviton, however, is still massless and thus the degeneracy in the gravitational supermultiplet is broken.

The breaking happens in the hidden sector with a vacuum expectation value $\langle F \rangle$ and a non-renormalizable coupling to the visible sector of strength $\sim 1/M_{\text{Planck}}$ as mentioned above.

⁵Spontaneous breaking of global Supersymmetry would require non-zero vacuum expectation values of either the F or D auxiliary fields (see section 2.2.2).

Figure 2.6: *Evolution of the scalar and gaugino mass parameters from the GUT scale to the electroweak scale for a MSUGRA/CMSSM scenario ($m_0 = 200$ GeV, $m_{1/2} = 600$ GeV, $A_0 = -600$ GeV, $\tan\beta = 10$, $\text{sign}\mu > 0$) based on the renormalization group equations. Taken from reference [43].*

In the case of vanishing gravitational interactions $M_{\text{Planck}} \rightarrow \infty$ and vacuum expectation value $\langle F \rangle \rightarrow 0$, the soft breaking mass terms m_{soft} must also vanish. This leads to the following approximative formula

$$m_{\text{soft}} \sim \frac{\langle F \rangle}{M_{\text{Planck}}}, \quad (2.129)$$

which implies for $m_{\text{soft}} \lesssim 1$ TeV a scale of $\sqrt{\langle F \rangle} \approx 10^{11}$ GeV for the hidden sector. The F-Field here refers to the auxiliary field of section 2.2.2, which is related to the superpotential as in equation 2.106.

An attractive, but quite constrained scenario, is that of minimal Supergravity (MSUGRA) which is often also referred to as constrained MSSM (CMSSM). It implies universal soft breaking terms and thereby reduces greatly the number MSSM parameters at the unification scale to:

- m_0 , the universal scalar mass,
- $m_{1/2}$, the universal gaugino mass,
- A_0 , the universal trilinear Higgs-sfermion-sfermion couplings,
- $\tan\beta$, the ratio of the vacuum expectation values of the two Higgs-doublets,
- $\text{sign}\mu$, the sign of the Higgsino mass parameter.

All parameters of the MSSM at the electroweak scale can be obtained from these five GUT scale parameters by application of the renormalization group equations as is illustrated in Fig. 2.6

Figure 2.7: *Mass spectrum of a typical MSUGRA scenario with $m_0 = 1450$ GeV, $m_{1/2} = 300$ GeV, $A_0 = 0$, $\tan \beta = 10$, $\mu > 0$ generated with ISAJET 7.58 [57].*

for one particular MSUGRA/CMSSM scenario. The mass spectrum of a given model at the electroweak scale depends on the chosen GUT-scale parameters. Fig. 2.7 shows the spectrum for a typical MSUGRA/CMSSM scenario.

2.2.6 Production and Decay Processes at the LHC

Supersymmetric particles, if they exist at energies accessible by the Large Hadron Collider, can be produced via the electroweak and strong interactions. Within R-parity conserving scenarios, such as the MSSM, particles are produced in pairs. The dominant production processes at the LHC are those of the strong interactions via gluon-gluon and gluon-quark processes and to a smaller extent via quark-quark interactions. The possible Feynman graphs of such processes at leading order are shown in Fig. 2.8. The electroweak production processes lead to charginos, neutralinos, and sleptons in the final state via electroweak vector bosons in the s-channel, and t-channel squark exchange. The corresponding diagrams can be found e.g. in reference [43]. For most supersymmetric mass spectra, the strong production cross-sections dominate over those of the weak processes as can be seen from Fig. 2.9.

Figure 2.8: *Leading order Feynman diagrams for the strong production of SUSY particles at the LHC, with gluon-gluon, gluon-quark, and quark-quark initial states from top to bottom. Taken from [43].*

Squarks typically decay, if kinematically allowed, through the process $\tilde{q} \rightarrow q\tilde{g}$ or otherwise to charginos or neutralinos according to $\tilde{q} \rightarrow q^{(\prime)}\tilde{\chi}^{0/(\pm)}$. The direct decay to the lightest neutralino $\tilde{\chi}_1^0$ is kinematically preferred and can dominate for right handed squarks since the $\tilde{\chi}_1^0$ has a large admixture of the Bino. The left-handed squarks, in turn, may preferentially decay into heavier charginos and neutralinos due to the stronger couplings to the Wino. Couplings of squarks to gauginos with a large Higgsino admixture are usually suppressed except for third generation squarks where the Yukawa couplings are large due to the heavy quark masses. Decays are also possible through virtual quarks if not kinematically allowed on-shell.

Gluinos decay through the process $\tilde{g} \rightarrow q\tilde{q}$ or if the latter is kinematically forbidden via a virtual squark directly to charginos or neutralinos according to $\tilde{g} \rightarrow q\tilde{q}^{(\prime)}\tilde{\chi}^{0/(\pm)}$. If squarks are

Figure 2.9: *Next-to-leading order MSSM SUSY production cross-sections in proton-proton collisions as a function of the average final state mass [58].*

not kinematically accessible the decay $\tilde{g} \rightarrow g\chi^0$ is possible.

As discussed previously the charginos and neutralinos are mixtures of the electroweak gauginos and Higgsinos. Depending on the mass spectrum of the SUSY scenario considered, the following two-body decays are in principle possible [43]:

$$\begin{aligned} \tilde{\chi}_i^0 &\rightarrow Z\tilde{\chi}_j^0, W\tilde{\chi}_j^\pm, h^0\tilde{\chi}_j^0, \tilde{l}\tilde{l}, \nu\tilde{\nu}, A^0\tilde{\chi}_j^0, H^0\tilde{\chi}_j^0, H^\pm\tilde{\chi}_j^\mp, q\tilde{q} \\ \tilde{\chi}_i^\pm &\rightarrow W\tilde{\chi}_j^0, Z\tilde{\chi}_1^\pm, h^0\tilde{\chi}_1^\pm, \tilde{l}\tilde{\nu}, \nu\tilde{l}, A^0\tilde{\chi}_1^\pm, H^0\tilde{\chi}_1^\pm, H^\pm\tilde{\chi}_j^0, q\tilde{q}' \end{aligned} \quad (2.130)$$

If these decays are kinematically excluded, they are replaced again by 3-body decays with off-shell gauge bosons, which decay on to fermions.

Sleptons typically decay to leptons and charginos or neutralinos according to $\tilde{l} \rightarrow l\tilde{\chi}^0, \nu\tilde{\chi}^\pm$ and $\tilde{\nu} \rightarrow \nu\tilde{\chi}^0, l\tilde{\chi}^\pm$. As for the quarks one has to distinguish between left- and right-handed sleptons, which preferentially decay to Wino or Bino like gauginos respectively.

Due to the cascade-like nature of supersymmetric decays described above, the expected experimental signature at the LHC for most MSSM-like models consists of several jets, missing transverse energy from the undetected LSP, and possible other objects, such as for example leptons.

2.2.7 Status and Implications of Experimental Bounds

While no evidence for supersymmetric particles has been found to date, experimental efforts have helped to significantly constrain the allowed regions of supersymmetric parameter space, in particular in the context of the well-studied MSUGRA/CMSSM scenarios. Some of the most important experimental bounds and their implications are described below. Figure 2.10 visualizes these bounds in the plane of universal scalar and gaugino masses, for scenarios with $\mu > 0$, $A_0 = 0$, and $\tan\beta = 10$ (left) as well as $\tan\beta = 50$ (right).

Cosmological constraints. It has previously been mentioned that in R-parity conserving SUSY models, the lightest supersymmetric particle (LSP) exhibits all features of a weakly interacting massive particle (WIMP), which in turn constitutes a candidate for cold dark matter. For MSUGRA/CMSSM scenarios this LSP with WIMP-features is the lightest neutralino $\tilde{\chi}_1^0$.

According to cosmological models, WIMPs were created in the early universe, when the energy density was high enough for the production processes. Thereafter they remained in thermal equilibrium until temperatures dropped below $m_{\tilde{\chi}_1^0}$ and production was suppressed. At the same time increasing distances between the particles reduced annihilation rates until today's "freeze-out" relic density $\Omega_{\tilde{\chi}_1^0}$ at a temperature of ~ 2.7 K was reached. Precise measurements

Figure 2.10: Overview of experimental and theoretical constraints on MSUGRA/CMSSM parameter space in the plane of universal scalar and gaugino masses, for scenarios with $\mu > 0$, $A_0 = 0$, and $\tan\beta = 10$ (left) as well as $\tan\beta = 50$ (right). The green shaded area corresponds to WMAP cosmological constraints [50], the blue shaded area to limits from LEP2 chargino searches [59], the blue dashed line to LHCb $B_s \rightarrow \mu^+\mu^-$ constraints [60], the orange shaded area to the region favoured by a_μ measurements [61], and the black solid line to the bounds from direct searches at the LHC [62, 63]. The red lines indicate the amount of fine-tuning according to the definition in reference [64]. Taken from reference [64].

Figure 2.11: Possible $b \rightarrow s\gamma$ processes in the Standard Model (top left) and supersymmetric models (top right and bottom). The photon line may be attached in all possible ways. Taken from reference [51].

of this relic density have been performed by the WMAP collaboration [50].⁶ The results can be translated into constraints of supersymmetric parameter space by the following mechanism:

The relic density is proportional to the inverse expectation value of the thermally averaged inclusive annihilation and co-annihilation cross section of the LSP pair times its relative velocity $\Omega_{\tilde{\chi}_1^0} \propto 1 / \langle \sigma v \rangle$ (e.g. [66]). The cross-section σ depends highly on the particular supersymmetric model and its associated masses and couplings. To find the allowed regions of parameter space for a model with universal parameters at the GUT scale like MSUGRA/CMSSM, one first needs to evaluate the renormalization group equations to obtain the MSSM parameters in the weak regime. Then a SUSY-spectrum generator can be used to determine higher-order contributions to masses and couplings. These are then fed into calculations of the annihilation inclusive cross section σ to yield the corresponding neutralino density. The bounds resulting from the latest WMAP measurements correspond to the green shaded areas in Fig. 2.10.

Constraints from $b \rightarrow s\gamma$. Another restriction on the parameter space results from measurements of the flavour-changing process $b \rightarrow s\gamma$. In the SM this decay involves loops containing W-bosons and up-type quarks as shown in Fig. 2.11 (top left). In the MSSM additional diagrams with SUSY contributions are possible as shown in Fig. 2.11 (top right and bottom). Measurements of the corresponding inclusive branching ratio, i.e. $b \rightarrow X_s \gamma$ (see [67]) put constraints on the kinematically allowed regions of parameter space, when compared with precise SM calculations.

Constraints from $B_s \rightarrow \mu^+ \mu^-$. Within the Standard Model the flavour-violating process $B_s \rightarrow \mu^+ \mu^-$ can occur at one-loop level with a highly suppressed branching ratio of the order of 10^{-9} . In MSUGRA/CMSSM scenarios this branching ratio is enhanced at large values of $\tan \beta$ due to additional flavour violating processes at one-loop level emerging from the supersymmetric Higgs sector. However, the upper limit on the $B_s \rightarrow \mu^+ \mu^-$ branching ratio is constrained by a recent measurement of the LHCb collaboration [60] to approximately less than two times the Standard Model expectation, which leaves little room for SUSY contributions. The region where the branching ratio calculated in MSUGRA/CMSSM falls below the experimentally allowed range is indicated with a dashed blue line in Fig. 2.10 (right).

⁶Measurements are also expected from the Planck collaboration (e.g. [65]).

$$\mu = \frac{eg}{2mc} \mathbf{S}, \quad (2.131)$$

Collider constraints. Direct searches at high energy physics collider experiments provide the most stringent bounds on the masses of supersymmetric particles. The current exclusion limits on gluino and squark production within MSUGRA/CMSSM based on LHC data taken in 2011 at 7 TeV center-of-mass energy are denoted by the solid black line labelled LHC7 in Fig. 2.10. The original results by the ATLAS and CMS collaborations are shown in Fig. 2.12.

Apart from the limits on sparticle masses from direct searches the most far-reaching implications on supersymmetric theories to date come from the recent observation of a new Higgs-boson like particle at the LHC. While in the Standard Model the mass of the Higgs boson is a free parameter of the theory as described in section 2.1.3, in the MSSM it is bound to values below the Z-boson mass in addition to radiative corrections, which lead to a total upper limit of $M_H \lesssim 135$ GeV [70–72]. The fact that the newly found particle lies within that rather narrow mass range may be interpreted as a hint in favour of weak scale Supersymmetry. However,

the measured mass range of this particle around 126 GeV requires a considerable amount of radiative corrections which are expected to originate mostly from the supersymmetric top quark. These corrections, in turn, require a careful tuning of the mass parameters of the MSSM to cancel contributions to the Higgs potential and to arrive at a vacuum expectation value within the electroweak regime. This problem is also known as the “little hierarchy problem”. The amount of fine-tuning⁷ that is necessary for a given MSUGRA/CMSSM scenario is indicated by the red lines in Fig. 2.10.

A number of theoretical proposals exist to alleviate the effects of a heavy Higgs boson on naturalness considerations. These are subject to ongoing studies and discussions inside the theoretical particle physics community. Two examples are given below.

Amongst one of the most well-known approaches is the so-called Next-to-minimal Supersymmetric Standard Model (NMSSM), which introduces a new gauge-singlet chiral supermultiplet S . In the simplest version of the NMSSM, the term $\mu H_u H_d$ in the MSSM superpotential (equation 2.126) is replaced by $\lambda H_1 H_2 S + \frac{\kappa}{3} S^3$ [73], where λ is the coupling of S to H_u and H_d , and κ is the self-coupling of the singlet field. In the NMSSM the μ -term is then generated dynamically via electroweak symmetry breaking where S takes the vacuum expectation value v_S resulting into an effective μ -term, $\mu_{eff} = \lambda v_S$. This mechanism allows to circumvent the so-called μ -problem of the MSSM [74], which emerges from the explicit appearance of the μ -term in the MSSM superpotential. While the value of μ is expected to be at the electroweak scale to allow for spontaneous symmetry breaking in the supersymmetric Higgs sector, there is no natural explanation why this value should be so small compared to e.g. the Planck scale, and why it should be of the same order of magnitude as the soft Supersymmetry breaking mass terms given that their physical origins are essentially unrelated. In the NMSSM the additional coupling λ of S to H_u and H_d can lead to larger masses of the lightest neutral CP-even Higgs boson than in the MSSM and thus makes the NMSSM a more natural candidate for Supersymmetry in the light of the recent experimental results.

Another path that is being followed by the theoretical community is that of “Natural” Supersymmetry (NSUSY). An overview can be found in [75] and references therein. NSUSY is a collective term for supersymmetric models, in which the squarks of the third generation are expected to be of the order of 1 TeV whereas the remaining supersymmetric quarks and leptons as well as the electroweak gauginos are essentially decoupled at energies of $\sim 10 - 50$ TeV. The gluino mass can be of the order of several TeV. These scenarios still allow for Supersymmetry with small fine-tuning and are to date consistent with the experimentally excluded limits on squark and gluino masses at the LHC.

⁷The high-scale fine-tuning definition described in reference [64] is used.

Chapter 3

Experimental Setup

This chapter gives an overview of the experimental setup used to perform the analyses described in this thesis. Section 3.1 outlines the main features of the Large Hadron Collider. Section 3.2 describes the ATLAS detector and its components.

3.1 The Large Hadron Collider

The Large Hadron Collider (LHC) [76] is currently the world's largest and most powerful particle accelerator. It is located at the European Organization for Nuclear Research (CERN) near Geneva in the Franco-Swiss border area. The LHC is designed to provide proton-proton collision with up to 14 TeV center-of-mass energy at an instantaneous luminosity of $10^{34} \text{ cm}^{-2}\text{s}^{-1}$. In addition the LHC can be operated with lead ions at 2.76 TeV center-of-mass energy. The collider ring is located 45-170 m underground in a circular tunnel that previously accommodated the LEP accelerator [77]. The LHC hosts seven experiments: ALICE [78], ATLAS [79], CMS [80], LHCb [81], LHCf [82], MoEDAL [83], and TOTEM [84]. These experiments are designed to exploit the full physics potential of the LHC with a diverse program of measurements and searches. Amongst the main motivations for the construction of the LHC are the searches for the elusive Higgs particle and for new physics signals beyond the Standard Model of particle physics.

In section 3.1.1 the full LHC acceleration chain is described. Section 3.1.2 gives an overview of important LHC beam parameters and their evolution during the first three data-taking periods.

3.1.1 LHC Accelerator Complex

Fig. 3.1 shows a sketch of the LHC accelerator setup complex. The protons used in the LHC are obtained via the ionization of hydrogen atoms using the Duoplasmatron proton ion source. From there the protons are passed to the Radio Frequency Quadrupole (RFQ), which organizes them into bunches and accelerates them over a distance of 1.75 m to a focused beam of 750 KeV. The RFQ is followed by the 30 m long Linac2 linear accelerator which further boosts the protons to 50 MeV. Then the protons are passed to the Proton Synchrotron Booster (PSB), a synchrotron consisting of four vertically stacked rings which accelerates the protons to 1.4 GeV. From the PSB the protons are injected in the Proton Synchrotron (PS) which has a

Figure 3.1: *Schematic drawing of the CERN accelerator complex. The protons bunches are passed through the Linac2, PSB, PS, and SPS facilities before being injected into the main LHC ring. The beams are progressively accelerated and eventually brought to collision in the LHC experiments. Figure not to scale. Adapted from reference [85].*

circumference of 628 m, and increases the beam energy to 26 GeV. The PS organizes the proton bunches into so-called bunch trains with approximately $1.15 \cdot 10^{11}$ protons per bunch and bunch spacing time intervals as low as to 25 ns at nominal running conditions. From the PS the bunch trains are injected into the Super Proton Synchrotron (SPS) which has a circumference of 6.9 km and increases the beam energy to 450 GeV. From the SPS the protons are finally injected via transfer lines into the LHC in clockwise and counter-clockwise directions. The LHC has a circumference of 26.7 km and accelerates the beam to its final energy using eight superconducting radio frequency cavities. The beams travel in separate vacuum beam pipes and are maintained in a fixed orbit using 1232 superconducting dipole magnets that are cooled to 1.9 K using liquid helium and provide a magnetic field of up to 8.33 T. In the transverse plane the beams are focused using 392 main superconducting quadrupole magnets. The maximum number of proton bunches circulating in the LHC is 2808. There are a number of empty time-slots between bunches or bunch trains which are necessary for the beam injection and abort procedure. Once the LHC beams have reached their final energy, they are declared as “stable” and the experiments located around the collision points on the LHC ring start collecting data for physics analysis. During a fill cycle the intensity of the beams gradually decreases due to collisions, beam-gas events, intra-beam scattering, and other effects. Once the intensity drops below a critical value, on average after 6 hours of stable beam operations, the beam reaches the end of its lifetime and is “dumped” by directing it against a target of absorbing material.

Parameter	2010	2011	2012	Nominal
Center-of-mass energy [TeV]	7	7	8	14
N_b [10^{11}]	1.2	1.5	1.6	1.15
$n_b^{\max}$	348	1380	1380	2808
ϵ_n [μm]	2.4-4	1.9-2.4	2.2-2.5	3.75
β^* [m]	3.5	1.5 $\rightarrow$ 1	0.6	0.55
Crossing angle [μrad]	200	240	290	285
Bunch Δt [ns]	150	75 $\rightarrow$ 50	50	25
μ	4	19	35	23
$L_{\max}^{\text{inst.}}$ [$10^{34} \text{ cm}^{-2}\text{s}^{-1}$]	0.02	0.4	0.76	1

Table 3.1: *Summary of important LHC beam parameters during the 2010, 2011, and 2012 proton-proton data-taking campaigns, as well as the nominal values of the design configuration.*

3.1.2 LHC Beam Parameters

The instantaneous luminosity at the interaction points along the LHC is determined by the beam parameters and can be calculated according to [76]

$$L = \frac{N_b^2 n_b f_{\text{rev}} \gamma_r}{4\pi \epsilon_n \beta^*} F, \quad (3.1)$$

where N_b is the number of protons per bunch, n_b the number of bunches per beam, f_{rev} the revolution frequency, γ the relativistic gamma factor, ϵ_n the normalized transverse beam emittance¹, and β^* the beta-function² at the collision point. F is the so-called geometric luminosity reduction factor, which arises from the non-zero crossing angle of the beams at the interaction point. In the nominal configuration the crossing angle is 285 μrad which corresponds to a luminosity reduction factor of 0.84. The nominal revolution frequency of the LHC is 40.079 MHz. A number of important beam parameters are shown in Tab. 3.1 for the nominal configuration as well as the actual values during the 2010, 2011, and 2012 data taking campaigns in proton-proton collisions. It can be seen that all parameter values gradually improved over the past three data-taking periods, some of them even beyond their design specifications. Tab. 3.1 also shows the time intervals between neighbouring bunches and the number of interactions per bunch crossing μ at peak luminosity. The latter quantity specifies the amount of so-called in-time pile-up interactions, to be distinguished from out-of-time pile-up which describes the overlap of interactions from two consecutive bunch crossings in a particle detector due to its finite readout timing resolution. The number of in-time pile-up interaction is proportional to the instantaneous luminosity and its constituting beam parameters.

Fig. 3.2 shows the evolution of the peak luminosity (top), the maximum mean number of events per beam crossing (middle), and the number of colliding proton bunches (bottom) as delivered to the ATLAS experiment by the LHC as a function of time over the 2010, 2011, and 2012 data-taking periods. Fig. 3.3 shows the integrated luminosity delivered to ATLAS to date in

¹The beam emittance is a measure of the average spread of the beam in position and momentum space.

²The value of β^* can be defined as the distance from the interaction point in the direction of the beam where the width of the beam envelope is twice as large as at the interaction point. It is a measure of how much the beam is squeezed towards the interaction point.

Figure 3.2: *Peak instantaneous luminosity (top), maximum mean number of events per beam crossing (middle), and the number of colliding proton bunches (bottom) as delivered to the ATLAS experiment by the LHC as a function of time over the 2010, 2011, and 2012 data-taking periods. Taken from [86].*

Figure 3.3: *Integrated luminosity delivered to ATLAS as a function of time over the 2010, 2011, and 2012 data-taking periods. Taken from [86].*

proton-proton collisions. It amounts to 48.1 pb⁻¹ and 5.61 fb⁻¹ at 7 TeV collision energy in 2010 and 2011 respectively and 23.3 fb⁻¹ at 8 TeV in 2012.

3.2 The ATLAS Detector

The ATLAS (A Toroidal LHC ApparatuS) experiment is a large general-purpose experiment located in an underground cavern at the LHC. The main objectives of ATLAS are precision measurements of the Standard Model of particle physics and searches for the Higgs boson as well as new physics beyond the Standard Model. The ATLAS detector has a forward-backward symmetric cylindrical shape with a length of 44 m, a diameter of 25 m, and a weight of approximately 7000 tons. ATLAS consists of several sub-detectors which are positioned in layers around the interaction point. Each of the main sub-detectors consists of so-called barrel elements which are aligned concentrically around the beam line and perpendicular endcap disks to cover the forward region.

Fig. 3.4 shows a computer-generated cut-away view of the ATLAS detector. Moving outwards from the interaction point the main components are the Inner Detector (ID), composed of the Silicon Pixel Detector, the Semiconductor Tracker (SCT), and the Transition Radiation Tracker (TRT), followed by the electromagnetic and hadronic calorimeters, and the muon spectrometer (MS). Separate solenoid and toroid magnet systems provide the magnetic fields for momentum reconstruction in the ID and MS respectively.

For the description of the ATLAS detector a right-handed coordinate system is used. The origin is located at the nominal interaction point, the positive x -axis points towards the center of the LHC, the y -axis upwards, and the z -axis into the direction of the beam. The parts of the detector at positive and negative values of z are referred to as “A-side” and “C-side”

Figure 3.4: Computer-generated cut-away view of the ATLAS detector indicating its main components. Taken from [79].

respectively. Spherical coordinates are defined by the radius R from the origin, the azimuthal angle ϕ in the transverse plane x - y with respect to the positive x -axis, and the polar angle θ with respect to the positive z -axis. The latter is often expressed in terms of pseudorapidity

$$\eta = -\ln \tan \frac{\theta}{2}, \quad (3.2)$$

which in the case of massless particles is equivalent to the rapidity

$$y = \frac{1}{2} \ln \frac{E + p_z}{E - p_z}, \quad (3.3)$$

of a particle with energy E and longitudinal momentum p_z . Using the η and ϕ coordinates the geometrical distance between objects in the detector can be expressed in terms of

$$\Delta R = \sqrt{\Delta \eta^2 + \Delta \phi^2}. \quad (3.4)$$

Pseudorapidity is the preferred representation of the polar angle, since differences in this quantity are Lorentz-invariant under a boost along in the z -axis. This is of particular relevance in hadron colliders like the LHC, where the momentum fraction of the interacting parton inside the hadron along the z -axis is not known.

Figure 3.5: Computer-generated cut-away view of the ATLAS inner detector indicating its main components. Taken from [79].

3.2.1 Inner Detector

The ID consists of the pixel detector close to the beam pipe, the SCT at intermediate radii, and the TRT at outer radii. The ID has a length of 6.2 m and a diameter of 2.1 m. The layout of the ID and its sub-detectors is illustrated in Fig. 3.5. The ID is designed to provide excellent pattern recognition and momentum resolution of charged-particle tracks as well as precise reconstruction of primary and secondary interaction vertices within $|\eta| < 2.5$. In addition, the TRT provides discrimination between electrons and hadron fakes within $|\eta| < 2.0$ using transition radiation. For the reconstruction of the track momentum and the identification of the sign of the particle charge the ID is immersed into a 2 T magnetic field, which is generated by a superconducting solenoid surrounding the ID. The solenoid has a length of 5 m and a diameter of 2.5 m. Its thickness amounts to only 45 mm corresponding to 0.66 radiation lengths³ which allows to minimize energy losses upstream of the calorimeters. The total material traversed by a particle in the ID in terms of radiation length X_0 and interaction length⁴ λ is shown in Fig. 3.6.

³A radiation length X_0 corresponds to the mean distance over which a high-energy electron loses the fraction $1 - (1/e)$ of its original energy by Bremsstrahlung processes where e denotes Euler's constant.

⁴The interaction length λ is defined as the mean free path of a high-energy particle in a material before undergoing inelastic scattering.

Figure 3.6: *Material distribution in terms of radiation length (left) and interaction length (right) as a function of pseudorapidity in the inner detector volume. Taken from [79].*

Silicon Pixel Detector

The silicon pixel detector consists of three barrel layers and three endcap disks inside a cylindrical envelope of 48.4 cm diameter and approximately 6.2 m length providing a pseudorapidity coverage of $|\eta| < 2.5$. It is instrumented with 1744 modules, each of which has a size of $19 \times 63 \text{ mm}^2$ and consists of 47232 silicon pixels with a minimum size of $50 \times 400 \text{ }\mu\text{m}^2$, corresponding to 46080 read-out channels in total. When a charged particle traverses the silicon, ionization processes create electron-hole pairs which are collected with an externally applied electric field and converted into a read-out signal. The accuracy of the pixel detector is $10 \text{ }\mu\text{m}$ in the R - ϕ plane and $115 \text{ }\mu\text{m}$ in z - and R -directions in the barrel and endcaps respectively. The innermost layer of the Pixel detector, referred to as the B-layer, is important for the reconstruction of secondary vertices which can arise from decays of long-lived particles such as B-hadrons, which may occur in close proximity of the primary interaction vertex. Being close to the interaction point the pixel detector and in particular the B-layer are subject to significant radiation doses. To reduce electronic noise related to radiation damage and other sources the detector is cooled to temperatures between -5 and $-10 \text{ }^\circ\text{C}$.

Semiconductor Tracker

The SCT consists of four barrel layers and nine endcap disks inside cylindrical volumes with a radius of 54.9 and 61.0 cm respectively. In total the SCT is equipped with 4088 modules resulting in approximately 6.3 million read-out channels. Each module consists of two pairs of single-sided silicon microstrip sensors which are glued back to back with an angle offset of 40 mrad to provide measurements in both coordinates of the plane. The microstrip sensors themselves consist of 70 - $90 \text{ }\mu\text{m}$ wide silicon strips. After the bonding of two sensors the combined strip length is 126 mm . The spatial resolution of the SCT amounts to $17 \text{ }\mu\text{m}$ in the R - ϕ plane and $580 \text{ }\mu\text{m}$ in z - and R -directions in the barrel and endcaps respectively. Straight tracks within $|\eta| < 2.5$ traverse at least four SCT modules resulting in four space-point measurements with two hits per module. Just like the pixel detector, the SCT is cooled to temperatures between -5 and $-10 \text{ }^\circ\text{C}$ to limit the effects of radiation damage.

Figure 3.7: *TRT high-threshold hit probability as a function pion and electron momentum in the barrel (left) and endcaps (right) respectively. Taken from [87].*

Transition Radiation Tracker

The TRT is built from approximately 351 000 gaseous straw tubes interleaved with transition radiation material. In the barrel section straws of 1.44 m length are aligned in parallel to the beam line in 73 layers between radii of 56.3 and 106.6 cm. The 37 cm long endcap straws are located in radial direction between 64.4 cm and 1.00 m, organized in 160 planes up to a distance of 2.7 m from the interaction point resulting in a full pseudorapidity-coverage of approximately $|\eta| < 2.0$. TRT tracks feature approximately 36 (22) hits when passing through the barrel section (endcaps).

The straw tubes of the TRT are 4 mm in diameter. They are filled with a xenon gas mixture (70 % Xe, 27 % CO₂, 3 % O₂) and a tungsten gold-plated wire located at the center. If a charged particle passes through the TRT it ionizes the gas inside the straw tubes. An electric field between the negatively charged straw tubes and the central wire induces to a localized ionization cascade. The ionization charges are collected at the central wire and transformed into a signal proportional to the energy of the original particle. From the drift times the original position of the particle inside the straw can be inferred. The single wire accuracy of a straw tube is approximately 130 μm and the maximum electron collection time amounts to 48 ns.

In addition, the tube layers of the TRT contain foils and fibers with different refractive indices which induce transition radiation photons if traversed by a charge particle. The energy of the transition radiation photons is proportional to the energy of the emitting particle divided by its mass. The response is therefore expected to be significantly higher for electrons than for hadronic particles. The transition radiation photons are absorbed by the Xenon gas giving rise to an ionization cascade. To distinguish between transition radiation and minimum-ionization signals low- and high-thresholds are used in the TRT front-end electronics. The discrimination power of the TRT between pions and electrons is illustrated in Fig. 3.7 which shows the high-threshold probability as a function of the pion and electron momentum.

Contrary to the pixel detector and the SCT, the TRT is operated at room temperature.

Figure 3.8: Computer-generated cut-away view of the ATLAS calorimeter system indicating its main components. Taken from [79].

3.2.2 Calorimeter System

The ID is surrounded by a calorimeter system which extends to 4.2 m in radius and 6.65 m along the positive and negative z -axis with a total coverage in pseudorapidity of $|\eta| < 4.9$. The purpose of the calorimeter system is to measure the energy and position of neutral and charged particles interacting with the calorimeter material. In the process of this measurement the particles are absorbed in the calorimeter material, generating a cascade of secondary particles with progressively degraded energies. Below a certain energy threshold ionization and excitation processes produce a detectable signal which manifests itself as charge or light emission. The energy of the initial particle can be reconstructed from these energy loss signals.

The nature of the particle cascades depends on the initial particle. Electrons and photons develop into electromagnetic showers via electron-positron pair production and the emission of photons via Bremsstrahlung. The showers terminate in the MeV regime where ionization, the photo-electric effect, and related processes start to dominate.

Hadrons initiate hadronic showers which are dominated by successive inelastic nuclear interactions. A significant fraction of the hadronic shower develops into electromagnetic showers due to the production of neutral pions and η -mesons which decay to photon pairs. In addition, some of the energy of the initial particle is absorbed in spallation and excitation processes and is therefore not detectable by the calorimeter (invisible energy). Further energy losses originate from the production of neutrinos or muons in the cascade which leave the calorimeter volume

undetected (escaped energy).

The average ratio of the calorimeter response to electromagnetic and hadronic particles of the same incident energy is used to classify a calorimeter as compensating or non-compensating. For compensating calorimeters this ratio is close to one, whereas for non-compensating calorimeters the response to electromagnetic showers is higher due to the invisible and escaped energy.

The ATLAS calorimeters are non-compensating and have a sampling geometry, where the energy deposition and signal detection are separated in alternating layers of absorbing passive material with high atomic number and a sensitive low density active detection medium. Sampling calorimeters offer the advantage of being relatively compact devices, as the passive layers can be used exclusively for energy dissipation. The possibility of fine segmentation in the longitudinal and lateral directions allows a measurement of the position with good accuracy. A drawback is the lack of precision in the energy measurement, which is due to fluctuations caused by the sampling structure.

The calorimeter energy resolution can be parameterized according to

$$\frac{\sigma_E}{E} = \frac{S}{\sqrt{E}} \oplus \frac{N}{E} \oplus C, \quad (3.5)$$

where $\oplus$ indicates the quadratic sum. The first term is referred to as stochastic term and accounts for statistical fluctuations related to the shower development, the sampling geometry, and the absorbing material in front of the calorimeter. The stochastic term has a smaller impact at higher energies where more particles are produced in the cascade and the impact of fluctuations is reduced. For sampling calorimeters the stochastic term is larger, due to variations in the number of observed particles in the active material. The second term in equation 3.5 is referred to as noise term and originates from electronic noise in the calorimeter read-out chain. Its impact on the energy resolution scales inversely with the energy and thus may affect the low-energy performance of calorimeters. The third term in equation 3.5 is the constant term which accounts for detector non-uniformities and calibration uncertainties. It limits the calorimeter performance at large energies.

The different sub-components of the ATLAS calorimeter system are shown in Fig. 3.8. Closest to the interaction point is the Liquid Argon (LAr) electromagnetic calorimeter (ECal) consisting of a barrel section (EMB) with coverage up to $|\eta| < 1.475$ and two end-cap wheels (EMEC) extending up to $|\eta| < 3.2$. The hadronic calorimeters are composed of scintillating tile calorimeters with one barrel and two extended barrel sections reaching up to $|\eta| < 1.7$, and two LAr hadronic end-caps (HEC) covering $|\eta| < 3.2$. In the forward region the LAr forward calorimeters (FCal) provide coverage within $3.1 < |\eta| < 4.9$. Details of the calorimeter sub-systems are given in the following paragraphs.

LAr Electromagnetic Calorimeter

The ECal consists of two symmetric half-barrel sections in positive and negative z -directions extending to $|\eta| = 1.475$, and two inner and outer end-cap wheels with coverage up to $|\eta| = 2.5$ and $|\eta| = 3.2$ respectively. The total calorimeter material traversed by electrons and photons corresponds to approximately 22 radiation lengths in the barrel and 24 radiation lengths in the endcaps. The ECal is a sampling calorimeter using LAr as active medium and lead as absorber

material. LAr was chosen due to its inherent radiation hardness and intrinsic linear behaviour with energy. A readout electrode consisting of kapton clad with copper is located at the center of each LAr layer. The sampling layers are arranged in an accordion like geometry as shown in Fig. 3.9 (left) which ensures full azimuthal coverage without dead regions and facilitates fast read-out at both ends of the electrodes. A nominal voltage of 2000 V is applied to the electrodes to collect the charge from ionization. The signals are processed by front-end crates located on the detector. Part of these crates are the front-end boards (FEBs) which perform analogue signal processing, digitization, and transmission to the read-out drivers (see section 4.1.1) located in a separate cavern in proximity of the detector.

The ECal is complemented by a presampler which covers regions $|\eta| < 1.52$ in the barrel and $1.5 < |\eta| < 1.8$ in the endcaps to correct for energy loss effects upstream of the calorimeter. The presampler is located in front of the ECal and contains one 11 mm (5 mm) thick active LAr layer in the barrel (endcaps) which are segmented according to $\Delta\eta \times \Delta\phi = 0.025 \times 0.1$. The ECal and presampler are enclosed in cryostat vessels which maintain a LAr temperature of 88 K.

As shown in Fig. 3.9 (left), the barrel component of the ECal consists of three longitudinal compartments with different segmentation. Within $|\eta| < 1.4$ the first compartment is finely granulated in η with strips of dimensions $\Delta\eta \times \Delta\phi = 0.0031 \times 0.1$ stretching over 4.3 radiation lengths. The excellent η -resolution allows to resolve and discriminate against close-by photons from pion decays. The middle sampling extends over 16 radiation lengths and consists within $|\eta| < 1.4$ of $\Delta\eta \times \Delta\phi = 0.025 \times 0.025$ segments absorbing the bulk of the shower energy. In the transition region $1.4 < |\eta| < 1.475$, which overlaps with the coverage of the endcaps, the first and second calorimeter layers are less finely instrumented. The third and last calorimeter layer covers regions $|\eta| < 1.35$ with a minimum thickness of two radiation lengths at $\eta = 0$ and a coarser granularity of $\Delta\eta \times \Delta\phi = 0.05 \times 0.025$.

The endcaps consist in the outer wheels of three layers within $1.5 < |\eta| < 2.5$ and two layers in the transition region $1.375 < |\eta| < 1.5$. The inner wheels are equipped with two layers which are coarsely granulated with segments of $\Delta\eta \times \Delta\phi = 0.1 \times 0.1$. The segmentation in the outer wheels is similar to that in the barrel and gets coarser with increasing pseudorapidity.

Tile Hadronic Calorimeter

The Tile hadronic calorimeter is located outside of the Ecal at radii between 2.28 - 4.28 m and consists of a barrel section covering $|\eta| < 1.0$ and two extended barrel sections within $0.8 < |\eta| < 1.7$. The Tile calorimeter is a sampling calorimeter using scintillating polystyrene tiles as active material and steel as absorber. The surface of the tiles is aligned in the R- ϕ plane as shown in Fig. 3.9 (right), providing almost complete azimuthal coverage. The ultraviolet light produced by ionizing particles passing through the scintillators is collected with wavelength-shifting optical fibers and transported to the outer end of the calorimeter where it is amplified by photo-multipliers and converted into an electronic signal. The fibers are grouped such that projective calorimeter towers in pseudorapidity are formed. The Tile calorimeter consists of three compartments. The first and second compartments have a cell granularity of $\Delta\eta \times \Delta\phi = 0.1 \times 0.1$ and the third layer is segmented according to $\Delta\eta \times \Delta\phi = 0.2 \times 0.1$. The material of the Tile calorimeter is equivalent to approximately 7.4 interaction lengths.

Figure 3.9: *Left: Illustration of the segmentation in the barrel component of the LAr electromagnetic calorimeter. Right: Schematic view of a wedge element of the Tile hadronic calorimeter. Taken from [79].*

LAr Hadronic Endcap Calorimeter

The LAr hadronic endcap calorimeter (HEC) covers the region $1.5 < |\eta| < 3.2$ in pseudorapidity and consist of two wheels on both sides of the detector built of 32 wedge-shaped elements each. The HEC is a sampling calorimeter using LAr as active medium which fills 8.5 mm wide gaps between absorbing copper plates of varying width. The ionization charges of traversing particles are collected with three electrodes which subdivide each gap into four separate 1.8 mm wide compartments. The granularity of the HEC is $\Delta\eta \times \Delta\phi = 0.1 \times 0.1$ in the region $1.5 < |\eta| < 2.5$ and $\Delta\eta \times \Delta\phi = 0.2 \times 0.2$ beyond. The HEC shares its cryostat vessel with the EMECs and the FCals. LAr is chosen as active medium in the forward region rather than scintillating tiles to profit from its inherent radiation hardness. The material of the HEC corresponds to approximately 10 interaction lengths.

LAr Forward Calorimeter

The LAr FCals are located between the beam line and the EMECs and HECs and further extend the pseudorapidity coverage up to $|\eta| = 4.5$. The FCals consists of three modules in longitudinal direction. The module closest to the interaction point is optimized for the measurement of electromagnetic objects and uses copper plates as the absorbing material. The other two modules use tungsten absorbers to obtain good containment of hadronic showers and limit punch-through of showers into the muon spectrometer. To handle the high particle flux close to the beam-pipe all modules feature fine LAr gaps for fast signal read-out. The cryostat structure is shared with the HECs and EMECs. The material of the FCals corresponds to approximately ten interaction lengths.

3.2.3 Muon Spectrometer

The MS forms the outermost part of the ATLAS detector. It extends over a range of 5-10 m in R , and 7-20 m in the positive and negative z -directions, enclosing a volume of 16 000 m³. The MS consists of a barrel section with three concentric cylindrical layers of muon stations within $|\eta| < 1$, and four endcap disks consisting of three main wheels located at 7 m, 13 m, and 21 m, as well as a smaller wheel at 11m, covering together the pseudorapidity region $1 < |\eta| < 2.7$. The MS is immersed in a magnetic field which provides bending power for muon momentum measurements. The magnetic field is generated by separate barrel and endcap superconducting toroid magnets which are enclosed in cryostat structures and provide field strengths up to 2.5 T. The MS consists of a number of different muon detection chambers. Monitored drift tube chambers (MDTs) and Cathode Strip Chambers (CSCs) are used for precision tracking. In addition there are resistive plate chambers (RPCs) and thin gap chambers (TGCs) for triggering and the measurement of the second coordinate in the non-bending plane. A cutaway view of the ATLAS muon system and its components is shown in Fig. 3.10.

The MDTs consist of two multi-layers of drift tubes. The drift tubes are made of aluminium and have a diameter of 3 cm. They are filled with a Ar/CO₂ gas mixture and a gold-plated tungsten wire is located at their center. Muons passing through the drift tubes create ionization charges which are collected by the anode wire to which a voltage of approximately 3 KV is applied. The drift tubes are positioned in the R - ϕ plane such that the center-points of the tubes are tangential to circles around the z -axis. In this way a resolution of 80 μ m per tube and 35 μ m per chamber is achieved in the bending plane. The MDTs are located in the barrel and endcap sections of the MS covering regions up to $|\eta| = 2.7$. Since MDTs can not handle large particle fluxes ($\gtrsim 150$ Hz/cm²), CSCs are used in immediate proximity of the beam line.

The CSCs are multi-wire proportional chambers consisting of four parallel planes. Each plane contains parallel anode wires which are aligned such that the central wire is pointing in the radial direction. The cathodes surrounding the wires are segmented with strips in parallel direction to the wires on one side and in perpendicular direction on the other. In this way each CSC plane provides position measurements in η and ϕ . Just like the MDTs the chambers are filled with a Ar/CO₂ gas mixture. The distance between the anode wires and between the anode wires and the cathode strips is 2.5 mm. The nominal voltage is 1.9 KV. The CSCs achieve a spatial resolution of 40 μ m in the bending plane and 5 mm in the non-bending direction. With electron drift-times of less than 40 ns the timing resolution per plane is approximately 7 ns. The CSCs are located approximately 7 m from the interaction point and approximately 90 cm from the beam line covering the pseudorapidity region $2.0 < |\eta| < 2.7$. They are inclined by 7 degrees towards the interaction point.

The RPC stations consist of two units with two independent detector layers each providing position measurements in η and ϕ respectively. Each RPC detector layer consists of two parallel resistive plates within a distance of 2 mm which are filled with an ionizing gas mixture (C₂H₂F₄/Iso-C₄H₁₀/SF₆). A voltage of 9.8 KV is applied between the plates. The signal from ionizing particles is read out with longitudinally and transversely segmented metal strips providing measurements of the η and ϕ coordinate respectively. The spatial accuracy of the RPCs is 10 mm. The timing resolution of 1.5 ns makes the RPC well suited for triggering. The pseudorapidity coverage of the RPCs installed in the MS is $|\eta| < 1.05$.

Similarly to the CSCs, the TGCs are multi-wire proportional chambers. The distance between the wires is 1.8 mm and the distance from the wires to the cathodes 1.4 mm. A voltage of 2.9

Figure 3.10: Computer-generated cut-away view of the ATLAS muon spectrometer and its main components. Taken from [79].

KV is applied. The interior of the chambers is filled with a $\text{CO}_2/\text{n-C}_5\text{H}_{12}$ highly-quenching gas mixture. The coordinate in the bending plane is obtained from ionization charges collected on the wires whereas the azimuthal coordinate is determined from strips perpendicular to the wires on the cathode side. The spatial accuracy of the TGCs is 2-6 mm in R and 3-7 mm in ϕ . The high voltage and the small distances between wires and cathodes result in a timing resolution of 4 ns. The pseudorapidity coverage of the TGCs installed in the MS is $1.05 < |\eta| < 2.7$ (2.4 for triggering).

3.2.4 Forward Detectors

ATLAS is equipped with a number of detectors in the forward region at large pseudorapidities. These are mostly used for luminosity measurements, the monitoring of LHC beam conditions, and the triggering of low- p_T scattering processes. Fig. 3.11 shows some of the detectors in the forward direction. With the exception of the FCal which has been treated in section 3.2.2 these detectors are described below.

Moving outwards from the interaction point, the first forward detector is the beam conditions

Figure 3.11: *Schematic view of the ATLAS forward detectors.*

monitor (BCM). The BCM monitors beam anomalies that could result into detector damage and performs measurements of bunch-by-bunch luminosities. It consists of four modules at a distance of 1.84 m from the interaction point in positive and negative z -direction. Each module contains two $500\ \mu\text{m}$ thick diamond sensors of size $1\ \text{cm} \times 1\ \text{cm}$ which are located 5.5 cm from the beam line and withstand large radiation doses. The sensors operate at a voltage of 1 kV to collect ionization charges of traversing particles. The low read-out latency at the order of ns allows a fast response in case of beam anomalies.

Next along the z -axis in positive and negative directions are the minimum bias trigger scintillators (MBTS) located 3.6 m from the interaction point. The MBTS consists of 32 scintillator counters which are installed on the inner cryostat walls of the endcap calorimeters between radii of 153-890 mm covering a region of $2.09 < |\eta| < 3.84$ in pseudorapidity. The light from the scintillators is transported with wavelength-shifting optical fibers to photomultipliers where it is amplified and converted into electronic signals. The MBTS was used for the triggering of low- p_T scattering events during the first days of LHC operation. Its usefulness has since decreased due to the expected radiation damage.

At $\pm 17\ \text{m}$ from the interaction point the luminosity monitor LUCID (Luminosity measurement Using a Čerenkov Integrating Detector) is located. It is used for relative luminosity measurement in ATLAS and covers regions of $5.6 < |\eta| < 6$ in pseudorapidity. LUCID consists of an array of 20 aluminium tubes of 1.5 m length and 15 mm diameter filled with C_4F_{10} gas. Charged

particles passing through the tubes emit Čerenkov light which is collected via photomultipliers at the end of the tubes. The timing resolution of LUCID is of the order of a few ns and thus allows the monitoring of individual LHC bunches.

The zero degree calorimeters (ZDCs) are located at ± 140 m along the z -axis, just behind the point where the ATLAS beam pipe branches into the two separate LHC beam pipes containing the counter-circulating proton beams. Due to this position the ZDCs can measure neutral particles, such as neutrons or photons in regions $|\eta| > 8.3$. The ZDCs provide an additional trigger for low- p_T scattering processes and allow to reduce backgrounds from interactions of protons with residual gas inside the beam pipe (“beam-gas events”) or with the beam collimators (“beam-halo events”). In heavy ion collisions ZDC information is important to determine the centrality of the events. The ZDCs are comprised of one electromagnetic and a number of hadronic modules, each of which consists of tungsten and steel absorber plates interleaved with quartz strips. The Čerenkov light created by the particle shower products in the modules is collected with photo multipliers.

The beam pick-up detectors (BPTX) are located at ± 175 m from the interaction point (not shown in Fig. 3.11). They are LHC beam position monitors which generate a high-precision timing signal for bunches passing through ATLAS. The BPTX signals are used as input to the ATLAS trigger system (see section 4.1) as well as for the monitoring of beam conditions and the timing of collisions with respect to the LHC clock.

At 237 m from the interaction point in positive and negative z -directions are the ALFA (Absolute Luminosity For ATLAS) forward detectors covering regions of $10.6 < |\eta| < 13.5$ in pseudorapidity. The purpose of ALFA is the absolute determination of the luminosity via measurements of the elastic proton-proton cross-section at small angles. ALFA consists of two so-called roman pot stations on each side of ATLAS. Each station contains two detector modules, one above and one below the beam. The modules consist of ten layers of scintillating fibers which provide a resolution of $30 \mu\text{m}$. ALFA is operated in special LHC runs, for which the detector modules can be approached up to 1 mm close to the beams.

3.2.5 Overview of Detector Performance

The performance goals of the main ATLAS detector components in terms of energy and momentum resolution are summarized in Tab. 3.2. Fig. 3.12 gives an overview of the resolutions actually achieved in reconstructed physics objects during operation. Fig. 3.12 (top) shows the

Detector component	Required resolution	$ \eta $ -coverage	
		Measurement	Trigger
Tracking	$\sigma_{p_T}/p_T = 0.05 \% p_T \oplus 1 \%$	2.5	-
Electromagnetic Calorimeter	$\sigma_E/E = 10\%/\sqrt{E} \oplus 0.7\%$	3.2	2.5
Hadronic Calorimeter	$\sigma_E/E = 50\%/\sqrt{E} \oplus 3\%$	3.2	3.2
Forward Calorimeter	$\sigma_E/E = 100\%/\sqrt{E} \oplus 10\%$	3.1 – 4.9	3.1 – 4.9
Muon Spectrometer	$\sigma_{p_T}/p_T = 10\% \text{ at } p_T = 1 \text{ TeV}$	2.7	2.4

Table 3.2: Resolution performance goals of the components of the ATLAS detector. Taken from [79].

Figure 3.12: *Top: Di-muon mass resolution for muons from Z-boson decays as a function of pseudorapidity in the mass range $75 \text{ GeV} < m_{\mu\mu} < 105 \text{ GeV}$ using the momentum measurements from the ID (left) and the MS (right) in 2012 ATLAS data. Bottom left: Di-electron invariant mass from Z-boson decays within the pseudorapidity region $|\eta| < 2.47$ using 2011 ATLAS data. Bottom right: Fractional jet energy resolution as a function of the average jet transverse momentum using 2011 ATLAS data. Taken from references [88–90].*

di-muon mass resolution for muons from Z-boson decays as a function of pseudorapidity in the mass range $75 \text{ GeV} < m_{\mu\mu} < 105 \text{ GeV}$ using the momentum measurements from the ID (left) and the MS (right) in 2012 ATLAS data. Resolutions of approximately 1.8 GeV in the ID and approximately 2.6 GeV in the MS are achieved depending on the pseudorapidity region. Figure 3.12 (bottom left) shows the di-electron invariant mass from Z-boson decays within the pseudorapidity region $|\eta| < 2.47$. The width of the distribution indicates a mass resolution of approximately 2 GeV. Finally, the fractional jet energy resolution as a function of the average jet transverse momentum is shown in Fig. 3.12 (bottom right), determined with in-situ techniques for a particular jet energy calibration scheme (“EM+JES”) in 2011 ATLAS data.

Further details on the detector performance are given in the context of the analyses presented in chapters 5 and 6.

Chapter 4

Data Acquisition and Data Processing

This chapter describes the acquisition and processing of data from the ATLAS detector. Section 4.1 reviews the online part of the processing system, in which the events of interest are selected and processed in real time as the particle collisions occur. The events are then written to storage where they are picked up by the offline data processing system, which prepares them for analysis and propagates them to ATLAS institutions world-wide. The offline data-processing system is described in section 4.2, including a short description of the Trigger - Tier0 liason task that the author fulfilled during the first two ATLAS data-taking periods. During offline processing the reconstructed physics objects and quantities are derived. The corresponding reconstruction algorithms of the objects and quantities relevant for this thesis are outlined in section 4.3. Finally, the simulation of events in ATLAS is described in section 4.4. The information given in this chapter is taken from the referenced sources.

4.1 Online Data-Processing

This section describes the main aspects of the online data-processing performed by the ATLAS trigger and data acquisition (TDAQ) system. Section 4.1.1 gives an overview of the architecture of the TDAQ system and describes the flow of data. Sections 4.1.2 - 4.1.4 describe the implementation of the main physics triggers used in the analyses presented in chapters 5 and 6.

4.1.1 The ATLAS Trigger and Data Acquisition System

During the 2011 data-taking campaign at a center-of-mass energy of 7 TeV the peak luminosity provided by the LHC reached a value of $3.65 \cdot 10^{33} \text{ cm}^{-2}\text{s}^{-1}$ with a maximum mean of approximately 19 interactions per bunch crossing with 50 ns between colliding bunches. At the design parameters of the LHC with 14 TeV center-of-mass energy and a luminosity of $10^{34} \text{ cm}^{-2}\text{s}^{-1}$, on average 25 interactions per bunch crossing at an event rate of 40 MHz are expected. These design values are expected to be exceeded when LHC operation resumes in 2015. The challenging task of the TDAQ system [79] is to efficiently select and record the rare events of interest to physics analyses and to reject the bulk of low transverse momentum background processes.

In ATLAS the trigger event selection is performed in three steps. Potentially interesting events are initially selected by the hardware-based first level trigger and passed on to the high-level trigger (HLT) for further scrutiny. The HLT consists of the software-based second-level trigger (L2) and the Event Filter (EF), which progressively refine the trigger decision and reduce the event rate to an acceptable level for offline processing and permanent storage.

A schematic diagram of the ATLAS TDAQ system is shown in Fig. 4.1. The first trigger level is implemented in custom-made electronics which allows fast processing times at the required event rate of 40 MHz. The L1 triggers use coarse granularity information from the electromagnetic and hadronic calorimeters as well as dedicated muon trigger chambers (RPCs and TGCs) to identify so-called regions of interest (RoIs) in terms of (η, ϕ) -coordinates in the detector. In addition a number of specialized triggers including the Minimum Bias Trigger Scintillators (MBTS) can be used at L1. No ID information is used at this stage since track reconstruction is not feasible in the available time intervals. The L1 trigger decision is made by the Central Trigger Processor (CTP), which combines the information of the L1 calorimeter and muon triggers on the multiplicity and energy thresholds of the coarsely reconstructed objects. The design rate of events accepted by L1 is 75 kHz. During the 2011 data taking campaign this value was around 65 kHz. The maximum latency to form the L1 trigger decision is $2.5 \mu\text{s}$. During this time the detector readout information for a given event is stored in pipeline memories, i.e. custom made front-end electronics, which are located either directly on or in close proximity of the detectors. If an event is accepted by L1, the data are transferred from the front-end electronics via Read Out Drivers (RODs) to Read Out Buffers (ROBs) for temporary storage. The RoIs are passed on to the software-based second level trigger which as part of the HLT runs on large server farms at CERN. At L2 more sophisticated reconstruction algorithms are executed which use in most cases full precision detector information inside the RoIs, including tracking information from the ID. The L2 trigger requests the necessary detector data inside these RoIs from the ROBs via the Read Out System (ROS). L2 has an average processing time of approximately 40 ms and is designed to reduce the event rate below 3.5 kHz. During the 2011 data-taking the maximum rate was typically 5-6 kHz. Events that are accepted by L2 are assembled by Subfarm Input nodes (SFIs) which collect full granularity detector information from the ROSs and build the entire event data structure to be further processed by the third level of the trigger system, the EF. Contrary to the previous trigger levels, the EF can use the full detector information provided in the event building process and runs complex reconstruction algorithms similar to those executed during offline data reconstruction (see section 4.3). The average latency at the EF is 4 s with a design output rate ranging between 200-300 Hz. The maximum output rate during the 2011 data-taking reached up to approximately 600 Hz. The events accepted by the EF are organized into several streams according to the trigger selection algorithms and written to the local CASTOR mass storage system at CERN by multiple instances of Subfarm Output (SFO) applications running on servers of the so-called data-logging computing farm. In case of transmission problems the data can also be buffered locally for up to approximately 24 h.

The output streams contain the relevant trigger and detector information in a raw data format and are subdivided into calibration stream, express stream, debug stream, and several inclusive physics streams. The calibration stream contains partial event information which is used for detector and trigger performance studies. The debug stream collects events with fatal or time-out processing errors which can be recovered at a later stage and help to understand and solve problems in the TDAQ system. The physics streams contain the complete events used for physics analyses and are subdivided according to trigger type into **Egamma**, **Muon**, **JetTauEtmis**, and **MinBias** streams. The express stream is formed as a small subset of the physics streams and

Figure 4.1: Schematic diagram of the flow of data and trigger information between the components of the ATLAS trigger and data acquisition system [91].

is used for offline data quality monitoring and the calculation of detector calibration constants which are required for the offline data reconstruction. The ensemble of triggers used for data-taking, so-called trigger menus, are defined in the trigger configuration database. Each trigger consists of a unique combination of L1 item and trigger chains at L2 and EF. The L1 item defines the selection criteria, in terms of energy thresholds or multiplicities of reconstructed objects at L1. The trigger chains define the sequence of steps to be followed at L2 and EF to determine the final trigger decision. The trigger chains consist of feature extraction algorithms (FEXs) for the reconstruction of trigger objects and quantities, and hypothesis algorithms (HYPOs), which apply selection criteria on the reconstructed objects. The execution of the steps inside a trigger chain is controlled by the trigger steering software [92].

The trigger menu also specifies the prescale factors which can be assigned to given L1 items or HLT chains to reduce their rate. The prescale factor p defines the fraction of events $1/p$ with positive trigger decision to be accepted at a given level. Individual events are selected

if a random number between 0 and 1 is smaller than $1/p$. Similarly pass-through factors can be defined which allow to accept events regardless of the trigger decision. The prescale or pass-through factors can be changed within a detector run to adjust the trigger rates.

The final objects and quantities associated to triggers are defined by the ATLAS event data model (EDM). They are created by the trigger algorithms at run time and are persisted and appended to the raw detector data in the output streams upon positive trigger decision.

The performance of the TDAQ and detector systems is monitored with a dedicated online monitoring framework, which analyses histogrammed information of quantities related to each subsystem. The histograms are automatically compared with reference histograms and green (“good”), yellow (“flawed”), or red (“bad”) status flags¹ are assigned. The flags are checked by data quality shifters. A final decision on the data quality is made after an offline evaluation as described in section 4.2.

4.1.2 Electron Triggers

Electron triggers are seeded by electromagnetic L1 items typically labelled $L1_EM[x]$, where $[x]$ refers to the transverse energy threshold of the corresponding L1 electromagnetic object in GeV measured at the precision of 1 GeV. At L1, RoIs are identified in the combined electromagnetic and hadronic calorimeters by forming so-called trigger towers with the reduced granularity of $\Delta\eta \times \Delta\phi \approx 0.1 \times 0.1$. A sliding window algorithm of 4×4 trigger towers scans the calorimeter volume within $|\eta| < 2.5$ in steps of 0.1 in η and ϕ . The horizontal or vertical energy sum of at least one of the four possible 2×1 or 1×2 combinations of neighbouring trigger towers in the 2×2 core of the 4×4 window is required to exceed a predefined threshold. In addition, isolation requirements are applied to the ring of cells around the 2×2 core in the electromagnetic calorimeter and to the ring and core cells in the hadronic calorimeter. The elements of the L1 cluster finding algorithm are illustrated in Fig. 4.2.

The η and ϕ coordinates of the RoIs which define the location of the 2×2 clusters identified at L1 are passed to the HLT. The trigger chains executed at L2 and EF are typically labelled $L2_e[x]_[y]$ and $EF_e[x]_[y]$, where e stands for electron, $[x]$ is the required transverse energy threshold in GeV, and $[y]$ classifies the electron identification criteria. Both at L2 and EF a sequence of FEX and HYPO algorithms is executed which form and refine electromagnetic clusters and ID tracks that are eventually combined to trigger electron objects. The cluster algorithm at L2 uses the second layer of the electromagnetic calorimeter at its full granularity of 0.025×0.025 in η and ϕ and determines a so-called pre-seed, defined as the cell with the largest energy deposition inside a given L1 RoI. The cluster position is then calculated from the energy weighted average cell positions in a 3×7 cell window around the pre-seed. As in the offline electron reconstruction algorithm described in section 4.3.1, the clusters are then formed by summing the energy depositions in 3×7 and 5×5 windows around the cell centers in the barrel ($|\eta| < 1.4$) and endcap ($1.4 < |\eta| < 2.47$) electromagnetic calorimeters respectively.

In a separate step tracks are built with fast tracking algorithms that fulfill the stringent timing requirements at L2. The tracking is performed around the L1 RoIs and uses hits from the silicon detectors. The reconstructed tracks are then extended to the transition radiation tracker if they fall into its acceptance. A minimum track- p_T of 1 GeV is required.

¹In addition, the colour grey indicates “undefined” usually due to insufficient statistics and the colour black is used if the corresponding subsystem is not operational.

Figure 4.2: *Illustration of the elements of the electromagnetic clustering algorithm at L1. Taken from [79].*

The electromagnetic clusters and tracks reconstructed at L2 are matched using $\Delta\eta_{\text{track-cluster}}$ and $\Delta\phi_{\text{track-cluster}}$ requirements. Further identification criteria on the shower shape variables R_η and E_{ratio} as well as hadronic leakage are applied. The definition of these variables will be discussed in section 4.3.1 and can be found in Tab. 4.2.

At the EF, sophisticated clustering and tracking algorithms are executed which are identical or very similar to those used in the offline electron reconstruction (see section 4.3.1). Consequently, the identification criteria applied to EF electron candidates resemble those of the offline reconstruction. Three sets of identification criteria with increasing background rejection power are defined. They are referred to as **loose**, **medium**, and **tight**. An additional scheme with tightened criteria labelled **loose1**, **medium1**, and **tight1** was created in 2011 to maintain low transverse energy trigger thresholds in the increased instantaneous luminosity and pile-up environment. The variables used for the different electron identification schemes at the EF are summarized in Tab. 4.1. Typically, the cut values on some of these variables are looser than those applied during offline electron reconstruction.

4.1.3 Muon Triggers

Muon triggers are seeded by L1 muon items typically labelled $\text{L1_MU}[\mathbf{x}]$, where MU stands for muon and $[\mathbf{x}]$ refers to the transverse momentum threshold of the corresponding L1 muon object measured in GeV. The L1 muon triggers are based on dedicated RPC and TGC trigger chambers (see section 3.2.3) in the barrel ($|\eta| < 1.05$) and endcaps ($1.05 < |\eta| < 2.4$) respectively. The L1 trigger algorithms look for coincidences of hits in different layers of chambers

Variable name	loose	medium	tight	loose1	medium1	tight1
η	•	•	•	•	•	•
$R_{\text{had}} \& R_{\text{had1}}$	•	•	•	•	•	•
R_{η}	•	•	•	•	•	•
$w_{\eta2}$	•	•	•	•	•	•
w_{stot}		•	•	•	•	•
E_{ratio}		•	•	•	•	•
n_{BL}			•		•	•
n_{pixel}		•	•	•	•	•
n_{si}		•	•	•	•	•
d_0		•	•	•	•	•
d_0^{tight}			•			
$\Delta\eta$		•	•	•		•
$\Delta\eta^{\text{tight}}$					•	
n_{TRT}			•			•
f_{TR}			•		•	•

Table 4.1: Criteria used in the different EF electron identification schemes. The definition of the variables can be found in Tab. 4.2. Table adapted from [93].

inside geometrical windows, so-called L1 trigger roads. The properties of these roads define the transverse momentum of the muon candidates at L1. Six thresholds can be configured, three at low transverse momentum (6-9 GeV), and three at high transverse momentum (9-35 GeV). The information of the barrel and endcap systems is collected in the CTP interface.

L2 uses the RoIs defined by L1 and runs a set of three algorithms. The first algorithm is a MS standalone algorithms that retrieves full granularity information from the MDTs inside the L1 RoIs and performs pattern recognition with subsequent track-fitting. In the next step, a fast combination algorithm combines the MS tracks with possible ID tracks to improve the resolution of the track parameters. Finally, an isolation algorithm is executed which examines energy depositions in the electromagnetic and hadronic calorimeters as well as ID tracks in the region of the muon candidate track.

At EF the full detector information from the the muon precision and trigger chambers is used. As at L2 a stand-alone algorithm is implemented which builds MS-only tracks and extrapolates them to the interaction point. In the so-called “outside-in” EF trigger the MS tracks are then combined with ID tracks. The complementary “inside-out” approach starts with ID tracks and combines them with MS tracks. The concrete implementation of EF algorithms is very similar to those used during offline muon reconstruction which is further described in section 4.3.2.

The HLT muon trigger chains are typically labelled $\text{L2_mu}[\mathbf{x}]$ and $\text{EF_mu}[\mathbf{x}]$, where mu stands for muon and $[\mathbf{x}]$ for the corresponding transverse momentum threshold.

4.1.4 Missing Transverse Energy Triggers

The missing transverse energy triggers are seeded by L1 missing transverse energy items labelled $\text{L1_XE}[\mathbf{x}]$, where XE stands for missing transverse energy and $[\mathbf{x}]$ for the corresponding threshold

at L1 in GeV. Since missing transverse energy is a global quantity, no RoIs are defined at L1. Instead the missing transverse energy is computed from the energy deposited in the same 0.1×0.1 trigger towers in η and ϕ as in the case of electron triggers (see section 4.1.2). All layers of the electromagnetic and hadronic calorimeters in a region of $|\eta| < 4.9$ are included. After pedestal subtraction and noise suppression the transverse energy in each tower is determined with look-up tables. The energies in blocks of four adjacent towers are then summed up vectorially to determine the x and y components of missing transverse energy. The L1 trigger decision is determined by comparing these values to pre-defined thresholds corresponding to the required value in missing transverse energy.

The absence of RoIs in the missing transverse energy triggers and the stringent timing constraints at L2 do not allow to access the entire calorimeter information at a finer granularity than at L1. Thus L2 is essentially based on the same detector information as L1. Instead of look-up tables as used at L1 for the determination of the missing transverse energy, the calculation at L2 is performed explicitly by taking the square root of the sum of the squared x and y components.

At the EF, the missing transverse energy is recalculated using the full precision calorimeter information and cuts are applied according to the required threshold.

At all trigger levels, the calorimeter energies are measured at the uncalibrated electromagnetic energy scale. No correction is applied for the lower calorimeter response to hadronic activity. Muons are not included into the missing transverse energy calculation. The performance of missing transverse energy triggers is thus evaluated by comparison to offline reconstructed missing transverse energy quantities where the muon contribution has been subtracted (see section 6.3).

The missing transverse energy trigger chains at L2 and EF are typically labelled `L2_xe[x]_noMu` and `EF_xe[x]_noMu`, where `xe` stands for missing transverse energy, `[x]` for the corresponding threshold in GeV, and `noMu` for the non-inclusion of muons into the calculation.

4.1.5 Jet Triggers

Jet triggers are seeded by L1 jet items typically labelled `L1_J[x]`, where `J` stands for jet and `[x]` refers to the corresponding transverse energy threshold in GeV. The basic building blocks of the L1 jet cluster reconstruction are jet elements consisting of 2×2 blocks of calorimeter trigger towers as defined in section 4.1.2. Within $|\eta| < 3.2$ jet RoIs are identified with a sliding window algorithm based on 4×4 , 6×6 , or 8×8 trigger tower windows. The energy depositions inside these windows, as well as the multiplicity of such elements, are compared to pre-defined values.

At L2 a simple cone algorithm² with three iterations is executed using the L1 RoI as seed. The L2 jet triggers were first activated at the end of the 2010 data-taking period.

At EF the entire event information is used to reconstruct jets with offline jet reconstruction algorithms as further described in section 4.3.3. With the beginning of the 2011 data-taking campaign the jet reconstruction at EF was performed unseeded, i.e. no geometrical information from L2 was used as input, but instead the entire event was scanned. EF jet triggers were first

²A cone algorithm uses the energy depositions inside a cone of given size around a pre-defined jet seed to recalculate the jet four momentum. The procedure is repeated iteratively until the jet position has stabilized.

actively used during the 2011 ATLAS data-taking. In the relatively low luminosity running conditions during 2010 there was no need for additional event rejection at the EF.

The jet trigger chains at L2 and EF with a multiplicity of one are typically labelled L2-j [x] and EF-j [x], where j stands for jet, and [x] for the corresponding transverse energy threshold in GeV.

4.2 Offline Data-Processing

After the acquisition of the data by the TDAQ system as described in section 4.1, the data are processed offline and prepared for analysis. Section 4.2.1 gives an overview of the data flow in the offline system from the initial processing at the CERN Tier-0 computing center to the world-wide distribution for user analysis. In this context section 4.2.2 gives a more specific summary of the ATLAS Trigger - Tier-0 liaison task that the author fulfilled during the first two ATLAS data-taking campaigns.

4.2.1 Data-flow in the Offline System

As described in section 4.1 the TDAQ system organizes the selected data events into several streams. The data are then picked up by the offline processing system for reconstruction and preparation for physics analysis by the CERN Tier-0 computing center. A schematic diagram of the data-flow at Tier-0 is shown in Fig. 4.3. Shortly after the beginning of a new data-taking run, Tier-0 rapidly processes a representative subset of the data contained in the express stream to derive detector calibrations and provide quick feedback on the data quality via offline monitoring tools. The detector calibrations, for example the beamspot position and dead or noisy channels of the sub-detectors, are uploaded to a detector conditions database. The data-quality is evaluated with automatic checks of histograms containing a variety of reconstruction-related quantities. As in the online monitoring system (see section 4.1), status flags according to a “traffic light” system are assigned to the histograms which are then inspected by data-quality shifters. The final “good” or “bad” data-quality decision is formed after combination of the online and offline flags involving if necessary the consultation of subsystem data-quality experts. The data quality decisions are later propagated to physics analyses in the form of so-called Good Run Lists (GRLs) which contain luminosity block³ ranges of detector runs that are deemed acceptable for a given physics analysis. 36 hours after the end of a data-taking run (“36 h calibration loop”) the bulk processing of the entire dataset from the physics streams is launched using the detector calibrations obtained from the express stream processing.

The reconstruction software derives the physics objects and quantities of interest as defined by the ATLAS EDM from the RAW input data. Along the offline processing chain, several output data formats based on the POOL persistency framework [95] with progressively smaller file sizes and reduced information are produced. The first and most detailed output format is the Event Summary Data (ESD), which contains detector-level hit information as well as reconstructed objects and quantities with an event size of approximately 1.5 MB. Analysis Object Data (AOD) is a subset of ESD and contains only the reconstructed objects and quantities relevant

³A luminosity block is a short time interval (typically 1-2 minutes) during a detector run in which the data-taking conditions (detector and beam conditions) were approximately stable.

Figure 4.3: Schematic diagram of the data flow during offline processing at the CERN Tier-0 computer center. Adapted from reference [94].

for physics analyses with an event size of approximately 150 KB. Finally the TAG format stores event-by-event metadata records including limited information on the most important reconstructed objects and quantities. The per-event size of the TAG format is approximately 4 KB. In addition several Derived Physics Data formats (DPD) exist. These are subsets of ESD or AOD datasets, where objects (“thinning”), parts of objects (“slimming”), or entire events (“skimming”) are removed. The per event size of DPDs thus depends on the exact selection applied. A particular type of DPD that is widely used in physics analyses is the D3PD format. Instead of the POOL persistency framework D3PDs are based on ROOT ntuples [96]. The D3PD content is tailored to specific classes of analyses.

With the exception of D3PDs⁴, the data formats described above are produced at the CERN Tier-0 computing center and subsequently distributed through the world-wide computing grid for further processing and storage. The ATLAS computing model foresees a hierarchical “Tier” structure of the computing facilities. According to this model the data are pushed from the Tier-0 computing center to ten Tier-1 centers distributed around the globe which are used for permanent storage, reprocessing campaigns, as well as the processing of jobs submitted through the computing grid interface. The tasks submitted via the grid interface can range from large scale Monte-Carlo production requests to user analysis jobs. Each Tier-1 is associated to a regional “cloud” of Tier-2 computing centers which are mainly used for data storage and processing of computing grid tasks. Each Tier-2 is associated to a number of local Tier-3 computing clusters. Tier-3s are mainly used for data analysis by local institutes, but can also

⁴D3PDs are centrally produced at Tier-1 and Tier-2 computing centers via the computing grid interface.

a release it was part of the author's responsibility to first evaluate the necessity of these updates at Tier-0. If adequately justified, the packages then underwent a set of reconstruction tests on cosmic ray, collision, and simulated datasets exercising the entire offline processing chain from RAW data files to ESDs and derived data formats as described above. If no problems were found, the software updates were added to a nightly validation release on which again a large suite of automated tests within the following testing frameworks was executed.

- The ATLAS Nightly Tests (ATN) are executed on the nightly built software releases and caches⁵. They contain a large number of tests checking various software aspects and run on input files with a few events only. The test output is automatically analysed and published on the nightly control system project page. Access to log-files is possible through the web interface.
- The Runtime Tester framework (RTT) is similar to ATN but runs on larger event samples and provides more tools for the analysis of the test output. It is thus possible to detect rare failures in the reconstruction software that did not emerge during ATN tests.
- The Tier-0 Chain Tests (TCT) run the full offline processing chain on a large data sample. TCTs automatically detect changes in the reconstruction output POOL files between two successive nightly releases. They are thus very useful to identify software packages which unintentionally change the output and therefore break the frozen Tier-0 policy (see below).

If any test failures were encountered the problem was identified and solved in collaboration with the responsible software developer. Upon successful completion of all tests, the release containing the new software packages was eventually deployed at Tier-0.

Due to the rapidly evolving running conditions and trigger menus during the 2010 and 2011 data-taking periods, the most common updates to the Tier-0 release were urgent improvements of the offline trigger monitoring. On a less frequent basis modifications to the EDM had to be introduced to Tier-0 to propagate new objects and quantities from the evolving online trigger software into the reconstructed data. Such updates had to adhere strictly to the so-called "frozen Tier-0 policy" which allows no changes to existing EDM objects in order to preserve consistency of the data within the same running period and to keep the evolving Tier-0 releases forwards and backwards compatible. EDM updates fulfilling these requirements - typically large bundles of inter-dependent software packages - were then jointly deployed online and at Tier-0.

Apart from handling new software requests, the Trigger - Tier-0 liason task also included the diagnosis of potential trigger-related problems in the existing Tier-0 release and their resolution in conjunction with the corresponding software developer.

The author fulfilled the Trigger - Tier-0 liason task from October 2009 throughout the first two ATLAS data-taking campaigns until the end of 2011.

4.3 Event Reconstruction

The reconstruction software running offline at the Tier-0 computing center derives the physics objects and quantities from the RAW input data which are then used for physics analysis. In

⁵A cache is a lightweight collection of a small number of software packages built on top of an existing release.

this section the reconstruction of electrons (section 4.3.1), muons (section 4.3.2), jets (section 4.3.3), missing transverse energy (section 4.3.4), and primary vertices (section 4.3.5) is described. These are the main physics objects and quantities used in the analyses presented in chapters 5 and 6. Details on the reconstruction of other objects including photons, tau leptons can be found e.g. in [79] and references therein.

4.3.1 Electron Reconstruction

The electron reconstruction in the central pseudorapidity region, $|\eta| < 2.47$, is based on electromagnetic calorimeter clusters which are associated to charged particle tracks reconstructed in the ID.

The basic building blocks for the reconstruction of electromagnetic clusters are longitudinal calorimeter towers of size 0.025×0.025 in η and ϕ corresponding to the granularity of the middle layer of the electromagnetic calorimeter. A sliding window algorithm identifies so-called seed clusters with a minimum transverse energy of 2.5 GeV using a window size of 3×5 calorimeter towers.

Tracks are identified in the ID within $|\eta| < 2.5$ using the standard ATLAS track reconstruction [97] based on the so-called inside-out⁶ and back-tracking⁷ algorithms. The tracks are extrapolated to the middle layer of the electromagnetic calorimeter and loosely matched to the seed clusters by requiring the angular distance between the extrapolated track and the electromagnetic cluster to satisfy $|\Delta\eta| < 0.05$ and $\Delta\phi < 0.1$ in the bending direction of the track to account for bremsstrahlung losses as well as $\Delta\phi < 0.05$ in the opposite direction. In the case of several tracks fulfilling these requirements, tracks with hits in the silicon detectors are preferred, and the one with the smallest value of $\Delta R = \sqrt{\Delta\eta^2 + \Delta\phi^2}$.

The electromagnetic clusters are then formed from the track-matched seed clusters using 3×7 and 5×5 blocks of longitudinal calorimeter towers in the barrel and endcaps respectively. The reconstructed electron energy is determined from the sum of the energy deposits in the cluster cells plus a number of corrections for the energy loss upstream of the calorimeter, as well as lateral and longitudinal energy depositions in and beyond the electromagnetic calorimeter (lateral and longitudinal leakage). The electron direction is determined from the η and ϕ directions of the reconstructed ID track at the vertex position.

Identification criteria are applied to reconstructed electrons to increase the fraction of true signal electrons and reject background electrons from photon conversions and pion decays as well as hadronic jets faking electrons. The variables used for electron identification are listed and described in Tab. 4.2. If the selection requirements are not explicitly specified in the table they are E_T - and η -dependent. In addition they may differ between the 2010 and 2011 versions of the electron identification to reflect the different pile-up conditions during the two data-taking campaigns and to improve the background rejection.

As for electron triggers, two schemes with three sets of progressively tighter identification criteria are defined. The nominal scheme, used both in 2010 and 2011, implements `loose`, `medium`, and `tight` selection criteria. The tightened and more pile-up robust “++”-scheme defines `loose++`, `medium++`, and `tight++` selections and is used only with data taken during the 2011 campaign.

⁶The inside-out tracking algorithm starts with three hits in the silicon detectors and subsequently includes more hit information at outer radii using a combinatorial Kalman filter. The tracks are extended to the TRT.

⁷The back-tracking algorithm uses TRT track segments and extends them to the silicon detectors.

Technically, the electron identification is implemented such that for a given reconstructed electron candidate binary status flags indicate which of the selection criteria from Tab. 4.2 are fulfilled. Based on a given pattern of status flags the electron candidate is then classified according to the identification schemes described above. A list of the selection requirements as a function of the identification scheme is given in Tab. 4.3. Note that, as described in Tab. 4.3 the actual cut values may vary between the nominal and “++”-schemes.

Type	Description	Name
Acceptance	$ \eta < 2.47$	
Hadronic leakage	Ratio of E_T in the first layer of the hadronic calorimeter to E_T of the EM cluster (used over the range $ \eta < 0.8$ and $ \eta > 1.37$)	R_{had1}
	Ratio of E_T in the hadronic calorimeter to E_T of the EM cluster (used over the range $ \eta > 0.8$ and $ \eta < 1.37$)	R_{had}
Middle layer of EM calorimeter	Ratio of the energy in 3×7 cells over the energy in 7×7 cells centred at the electron cluster position	R_η
	Lateral shower width, $\sqrt{(\sum E_i \eta_i^2)/(\sum E_i) - ((\sum E_i \eta_i)/(\sum E_i))^2}$ where E_i is the energy and η_i is the pseudorapidity of cell i and the sum is calculated within a window of 3×5 cells	$w_{\eta 2}$
Strip layer of EM calorimeter	Shower width, $\sqrt{(\sum E_i (i - i_{\text{max}})^2)(\sum E_i)}$, where i runs over all strips in a window of $\Delta\eta \times \Delta\phi \approx 0.0625 \times 0.2$, corresponding typically to 20 strips in η , and i_{max} is the index of the highest-energy strip	w_{stot}
	Ratio of the energy difference between the largest and second largest energy deposits in the cluster over the sum of these energies	E_{ratio}
Track quality	Number of hits in the pixel detector (≥ 1)	n_{pixel}
	Number of total hits in the pixel and SCT detectors (≥ 7)	n_{si}
	Transverse impact parameter ($ d_0 < 5$ mm)	d_0
	Tighter transverse impact parameter requirement ($ d_0 < 1$ mm)	d_0^{tight}
Track-cluster matching	$\Delta\eta$ between the cluster position in the strip layer and the extrapolated track ($ \Delta\eta < 0.01$)	$\Delta\eta$
	$\Delta\phi$ between the cluster position in the middle layer and the extrapolated track ($ \Delta\phi < 0.02$)	$\Delta\phi$
	Ratio of the cluster energy to the track momentum	E/p
	Tighter $\Delta\eta$ requirement ($ \Delta\eta < 0.005$)	$\Delta\eta^{\text{tight}}$
TRT	Total number of hits in the TRT	n_{TRT}
	Ratio of the number of high-threshold hits to the total number of hits in the TRT	f_{TR}
Conversions	Number of hits in the b-layer (≥ 1)	n_{BL}
	Veto electron candidates matched to reconstructed photon conversions	

Table 4.2: Summary of electron identification variables. The cut values depend on the underlying identification scheme and may be E_T - and η -dependent. Where this is not the case example cut values are shown which may not be applicable to all identification schemes. Table adapted from [98].

4.3.2 Muon Reconstruction

There are four distinct strategies for the reconstruction of muons in ATLAS [99]. Stand-alone muons are reconstructed in the MS only and are extrapolated back to the beam line. They are corrected for energy losses upstream of the MS. Combined muons are based on the combination of stand-alone muons reconstructed in the MS with tracks reconstructed in the ID. Segment-tagged muons start with a reconstructed track in the ID and look for matching track segments

Variable name	loose	medium	tight	loose++	medium++	tight++
$ \eta < 2.47$	•	•	•	•	•	•
$R_{\text{had}} \ \& \ R_{\text{had1}}$	•	•	•	•	•	•
R_η	•	•	•	•	•	•
$w_{\eta 2}$	•	•	•	•	•	•
w_{stot}		•	•	•	•	•
E_{ratio}		•	•	•	•	•
n_{BL}			•		•	•
n_{pixel}		•	•	•	•	•
n_{si}		•	•	•	•	•
d_0		•	•		•	•
d_0^{tight}			•			
$\Delta\eta$		•	•	•	•	•
$\Delta\eta^{\text{tight}}$			•			
$\Delta\phi$			•			•
E/p			•			•
n_{TRT}			•			•
f_{TR}			•		•	•
Conversion Match			•			•

Table 4.3: *Criteria used in the different electron identification schemes. The definition of the variables can be found in Tab. 4.2. The cut values may differ between the nominal and “++”-schemes such that e.g. the application of d_0^{tight} and $\Delta\eta^{\text{tight}}$ criteria is not necessary in the “++”-schemes.*

in the MS. Finally calorimeter-tagged muons associate ID tracks with energy depositions in the calorimeters which are compatible with the signature of minimum ionizing particles.

There are several algorithms that implement the reconstruction strategies described above. The reconstructed muons used in this thesis are based on the so-called **STACO** algorithm chain containing the algorithms **MuonBoy**, **MuTag**, and **STACO** itself [100]. The **MuonBoy** algorithm produces stand-alone muons starting from so-called regions of activity in η and ϕ identified by the RPCs and TGCs. Local track segments are reconstructed in these regions of activity in all muon stations. The track segments are combined to form track candidates followed by a global track fit throughout the entire MS to determine the final track parameters. The track is extrapolated to the beam line and corrections for energy loss in the calorimeters are applied. The **STACO** algorithm is based on the statistical combination of independently measured tracks in the MS and the ID which provides an improved momentum resolution and the possibility to reject muons from secondary interactions. Finally the **MuTag** algorithm tags those ID tracks not used by **STACO** as muons if they match a track segment in the MS. The **MuTag** algorithm allows to recover muons with low transverse momentum whose tracks may only be identifiable in the inner-most layers of the MS.

4.3.3 Jet Reconstruction

Jets can be defined at parton level, particle level, and detector level. At parton level the jet corresponds to the final state particles of the hard scattering process that develop into jets, whereas at particle level it consists of the stable particles after fragmentation and hadronization of the final state partons. At detector level the jets are calculated from the energy depositions of the stable particles inside the detector volume. While the concepts of parton and particle

level jets are relevant in the context of theoretical calculations and Monte-Carlo simulations, only detector level jets can be reconstructed in experimental data.

The basic building blocks of the jet reconstruction are either calorimeter towers or topological calorimeter clusters (short topoclusters). Calorimeter towers are, as previously defined, regions in η and ϕ in which the energy deposits in the calorimeter cells are summed up in radial direction. Further details on the exact specification of calorimeter towers used in jet reconstruction algorithms in ATLAS can be found in [101]. The reconstructed jets used in the analysis presented in chapter 6 are based on topoclusters. Topoclusters are three-dimensional objects of calorimeter cells. They are formed around seed calorimeter cells that have a large excess of energy compared to the expected noise level. In an iterative procedure neighbouring cells whose energy exceeds a lower threshold over the expected noise level are then grouped together with the seed cell. The topoclusters formed in this way are finalized by adding the directly neighbouring cells as outermost layer with no additional threshold requirements. Topoclusters containing significant local energy deposition maxima are split.

The jets used in this thesis are reconstructed with the **AntiKt** algorithm [102] based on topoclusters. The **AntiKt** algorithm belongs to the class of sequential clustering algorithms in which distance measures between two topoclusters i and j ,

$$d_{ij} = \min(k_{ti}^{2p}, k_{tj}^{2p}) \frac{\Delta_{ij}^2}{R^2}, \quad (4.1)$$

as well as between a topocluster i and the beam B ,

$$d_{iB} = k_{ti}^{2p}, \quad (4.2)$$

are defined. The transverse momentum of object i is denoted k_{ti} and

$$\Delta_{ij}^2 = \Delta y_{ij}^2 + \Delta \phi_{ij}^2 \quad (4.3)$$

is the radial distance between two objects with differences in rapidity and azimuthal angle Δy_{ij} and $\Delta \phi_{ij}$. R is the radius parameter of the algorithm. The clustering algorithm compiles a list of distances d_{ij} between the topoclusters in the event and merges those with the smallest distances. The list is remade and the process is repeated iteratively. If $d_{iB} = k_{ti}^{2p}$ is the smallest distance, the object is called a jet and is removed from the list. In the **AntiKt** algorithm the parameter p in equations 4.1 and 4.2 is set to a value of -1. The **AntiKt** algorithm is unaffected by the emission of soft particles or by the splitting of particles into two collinear particles. These are two important properties of the jet reconstruction referred to as infrared and collinear safety.

Without further corrections, the energy of the reconstructed jets is measured at the electromagnetic scale, i.e. the default energy scale at which the calorimeters are calibrated. To correct for the lower energy response in hadronic showers due to the non-compensating character of the calorimeters (see section 3.2.2), a p_T - and η -dependent calibration according to the so-called EM+JES scheme is applied. The EM+JES calibration is obtained from comparisons of reconstructed jets to truth jets in Monte-Carlo simulation including corrections for pile-up interactions, and the primary vertex position as further described in [103].

4.3.4 Missing Transverse Energy Reconstruction

The determination of missing transverse energy in the ATLAS detector is based on energy deposits in the calorimeters as well as reconstructed muons. The components of missing transverse energy in x and y direction are thus divided into a calorimeter and a muon term according to

$$E_{x(y)}^{\text{miss}} = E_{x(y)}^{\text{miss,calo}} + E_{x(y)}^{\text{miss},\mu}. \quad (4.4)$$

The calorimeter term can be determined either in a purely cell-based way summing up the energy deposits in the individual calorimeter cells, or in an object-based fashion associating the calorimeter cells to reconstructed physics objects and using their dedicated calibration schemes. The latter approach is adopted for the missing transverse energy quantity used in this thesis. In the analysis presented in Chapter 6, the corresponding missing transverse energy calorimeter term is calculated according to

$$E_{x(y)}^{\text{miss,Calo}} = E_{x(y)}^{\text{miss,e}} + E_{x(y)}^{\text{miss,jets}} + E_{x(y)}^{\text{miss,CellOut}}, \quad (4.5)$$

where each term is calculated from the calorimeter cells associated to the corresponding objects inside the pseudorapidity range $|\eta| < 4.5$ according to

$$E_x^{\text{miss,term}} = - \sum_{i=1}^{N_{\text{cells}}^{\text{term}}} E_i \sin \theta_i \cos \phi_i \quad \text{and} \quad E_y^{\text{miss,term}} = - \sum_{i=1}^{N_{\text{cells}}^{\text{term}}} E_i \sin \theta_i \sin \phi_i \quad (4.6)$$

for the x and y coordinates respectively. In the default configuration, the term $E_{x(y)}^{\text{miss,e}}$ is calculated from the cells inside electromagnetic clusters which are associated to **medium** electrons with transverse energy $E_T > 10$ GeV calibrated at the electron energy scale. The term $E_{x(y)}^{\text{miss,jets}}$ is determined from the topoclusters associated to the jets reconstructed as described in section 4.3.3 with $p_T > 20$ GeV calibrated at the jet energy scale. The term $E_{x(y)}^{\text{miss,CellOut}}$ contains all cells from topoclusters that do not belong to any of the reconstructed objects described above, calibrated at the jet energy scale.

The terms included in equation 4.5 and their exact definition can vary from analysis to analysis. Typically each term is recalculated at analysis level to propagate the exact object definitions to the missing transverse energy calculation (see section 6.4.4). Additional terms of τ -lepton, low- p_T jets and photons are sometimes included into the calculation if they are not absorbed in the $E_{x(y)}^{\text{miss,CellOut}}$ term.

In the default configuration the muon term is calculated from the transverse momentum of the muons in the event within $|\eta| < 2.7$ according to

$$E_{x(y)}^{\text{miss},\mu} = - \sum_{i=1}^{N_{\text{muons}}} p_{x(y)}^{\mu}. \quad (4.7)$$

The definition of muons included into the default calculation is given in [104]. In the analysis described in chapter 6 the muon term is recalculated with the exact muon definition used in the analysis.

The final missing transverse energy quantity and its azimuthal direction is determined according to

$$E_T^{\text{miss}} = \sqrt{(E_x^{\text{miss}})^2 + (E_y^{\text{miss}})^2} \quad \text{and} \quad \phi^{\text{miss}} = \arctan(E_y^{\text{miss}}, E_x^{\text{miss}}). \quad (4.8)$$

4.3.5 Primary Vertex Reconstruction

The position of primary interaction vertices is determined in two steps. Reconstructed tracks fulfilling a number of quality requirements on the transverse momentum, transverse and longitudinal impact parameters and their uncertainties, and the number of hits in the silicon detectors are used to establish primary vertex candidates. In a second step a vertex fitting algorithm determines the vertex position and its uncertainty inside the interaction region. Several primary vertex reconstruction algorithms implement these two steps using either the so-called “fitting-after-finding” approach in which the vertex fit is performed after the determination of the vertex position, or the “finding-through-fitting” method which determines the vertex position during the fitting procedure. Details on the primary vertex reconstruction can be found in references [105, 106].

4.4 Event Simulation

The production of simulated event samples in ATLAS proceeds in several steps. First, Monte-Carlo generator programs are used to model the fundamental physics processes of the proton-proton collisions and the subsequent particle decays as further described in section 4.4.1. The output of Monte-Carlo generation, typically in the form of particle four-momentum vectors, is passed to the detector simulation which describes the interaction of the final state particles with the detector material. The detector simulation is based on modeling of the passage of particles through matter implemented in the GEANT4 package [107] using a detailed description of the geometry of the ATLAS detector. The digitization step then transforms the output of the detector simulation into raw readout signals as they would be produced in a real event. Finally the samples are reconstructed with the same reconstruction algorithms as previously described in section 4.3 leading to the same output formats (see section 4.2) as for real data events. To preserve consistency very similar software releases are used for the processing of data and simulated samples.

The detector simulation is the computationally most expensive step in the event simulation chain. In view of finite computing resources and large Monte-Carlo production campaigns the full simulation with GEANT4 can be replaced with fast simulation using ATLFAST-II [108]. ATLFAST-II in its default configuration is based on the parameterization of particle showers in the calorimeters which in full simulation make up the bulk of the computing time. The parameterizations are tuned and validated with data and full simulation samples. In addition the fast simulation of tracks in the ID and MS is possible.

4.4.1 Monte-Carlo Generation

Monte-Carlo generators randomly sample from probability models of given particle physics processes to obtain simulated event samples. Due to the factorization theorem of QCD [109], the modeling of a particle collision and the production of the final state particles can be divided into several steps. These steps correspond to the perturbative and non-perturbative regimes of the strong interactions which are separated by a factorization scale μ_F . The separation is necessary since the strong coupling constant α_s evolves with increasing interaction energy to smaller values as predicted by the renormalization group equations 2.87. Fig. 4.5 illustrates the structure of a proton-proton collision event at the LHC. According to the quark-parton model each proton consists of three so-called valence quarks and a sea of gluons and quark-antiquark pairs whose effects are resolved at high scattering energies. The distribution of partons inside a proton is described by parton density functions (PDFs). When the two incoming protons undergo a deep inelastic scattering process, the corresponding cross-section is a convolution of PDF effects which are inherently non-perturbative and the hard interaction of two partons with large momentum transfer described by perturbative QCD calculations. The partons in the initial and final state of the hard interaction form so-called parton showers by radiating virtual gluons which emit further gluons or produce quark-antiquark pairs. If the energy of the constituents of the parton shower reaches the non-perturbative regime at low energies and thus high values of the strong coupling constant, the process of hadronization sets in. During hadronization the colour-charged partons develop into colour-neutral hadron states which may be unstable and decay further into the final state particles of the event. Along with the hard scattering process of interest, there is additional activity in the event associated to the remnants of the protons and possibly due to pile-up interactions from other proton-proton collisions occurring during the same bunch-crossing. Here such processes are collectively termed as the underlying event.

Hard Process

The cross-section for a hard process $pp \rightarrow n$ in a proton-proton collision can be calculated as convolution of the parton distribution functions $f_i(x_i, \mu_F)$ of the participating protons and the parton-level cross-section $\hat{\sigma}_{ab \rightarrow n}$ according to [111]

$$\sigma_{pp \rightarrow n} = \sum_{a,b} \int_0^1 dx_a dx_b \int f_a(x_a, \mu_F) f_b(x_b, \mu_F) d\hat{\sigma}_{ab \rightarrow n}(\mu_F, \mu_R). \quad (4.9)$$

The indices a and b run over the partons inside the proton and μ_F and μ_R correspond to the factorization and renormalization scales respectively. In a proton-proton interaction with momentum transfer at the scale μ_F , the parton distribution function $f_i(x_i, \mu_F)$ gives the probability to find parton i with momentum fraction $x_i = p_i/P$ inside a proton with momentum P . While the scale dependence of the PDFs can be described perturbatively by the Dokshitzer-Gribov-Lipatov-Altarelli-Parisi (DGLAP) formalism [112–114], the evolution with x falls into the non-perturbative regime and must therefore be extracted from fits to experimental data. Several groups, such as the CTEQ [115] and MSTW [116] collaborations, provide parameterized PDF sets using data from a number of fixed target (e.g. NuTeV [117]) and collider (e.g. HERA [118, 119], Tevatron, LHC) experiments. The parton level cross-section $\sigma_{ab \rightarrow n}$ for the production of final state n from partons a and b is calculated perturbatively from the transition matrix element squared $|\mathcal{M}_{ab \rightarrow n}|^2$ by summing over the Feynman diagrams contributing

Figure 4.5: *Illustration of the various steps of event generation. Underlying event as well as initial state radiation processes are not shown. Figure adapted from [110].*

to the process and averaging over the spin and colour degrees of freedom. The correct choice of the factorization and renormalization scales μ_F and μ_R is not known a priori and introduces some degree of arbitrariness into the cross-section calculation. The default choices are typically related to characteristic quantities of the corresponding processes such as e.g. the mass of a resonance [120]. Uncertainties related to this choice are usually evaluated with scale variations by a factor of two in positive and negative directions. Most of the existing Monte-Carlo generation tools provide the calculation of the parton level cross-section of $2 \rightarrow 2$ and $2 \rightarrow n$ processes at leading order (LO) in perturbation theory which provides to good accuracy a description of the shape of the kinematic variables of interest but often mismodels the absolute normalization. For this purpose so-called normalization K-factors can be applied which normalize the distributions to the corresponding cross-sections at next-to-leading order accuracy (NLO). There are also several Monte-Carlo generation tools which are fully implemented at NLO accuracy taking into account Feynman diagrams with additional virtual particle loop or real particle radiation corrections.

Parton Shower

The final state partons originating from the hard process are subject to QCD radiation. The successive emission of gluons and the production of quark-antiquark pairs leads to the development of parton showers, which represent higher order corrections to the hard process that cannot be calculated explicitly. The evolution of a parton shower is described by an evolution variable t , usually taken as the virtuality, i.e. the virtual mass squared of the parton, or the opening angle of the corresponding parton shower process. The matrix elements involving parton shower processes are strongly enhanced at low parton energies and small opening angles. In the limit of soft and collinear emission this results into divergencies which are avoided by terminating the shower at a cut-off scale t_0 .

In the case of near-collinear splitting of final state partons of flavour i into partons j and k the cross-section of the hard process at leading order can be factorized according to [111]

$$d\sigma \approx \sigma_0 \sum_{\text{partons}, i} \frac{\alpha_s}{2\pi} \frac{d\theta^2}{\theta^2} dz P_{ji}(z, \phi) d\phi, \quad (4.10)$$

where $P_{ji}(z, \phi)$ are splitting functions which depend on the flavour of the participating partons, on the energy fraction z of parton i carried by parton j , and on the azimuthal angle ϕ of the splitting. Further, σ_0 corresponds to the partonic cross-section of the hard interaction at matrix element level, θ to the opening angle of the splitting, and α_s denotes the strong coupling constant. There are four possible splitting functions corresponding to the processes $g \rightarrow qq$, $g \rightarrow gg$, $q \rightarrow qg$, $g \rightarrow gq$.

Based on equation 4.10 the parton shower development is implemented in an iterative procedure using the so-called Sudakov form factor [111]

$$\Delta(t_i^2, t_{i+1}^2) = \exp \left(- \int_{t_{i+1}^2}^{t_i^2} \frac{dq^2}{q^2} \frac{\alpha_s}{2\pi} \int_{t_0^2/q^2}^{1-t_0^2/q^2} dz P_{ji}(z) \right), \quad (4.11)$$

which corresponds to the non-splitting probability of parton i evolving from scale t_i to t_{i+1} . At each step of the iteration, equation 4.11 is solved for a random value of $\Delta(t_i^2, t_{i+1}^2)$ drawn from a uniform probability distribution between 0 and 1. The parton shower is terminated if the non-perturbative regime at the cut-off value $t_{i+1}^2 < t_0^2 \approx 1 \text{ GeV}^2$ is reached.

For the generation of initial state radiation parton showers, the PDFs of the colliding protons have to be included into the calculation of the Sudakov form factors. To avoid inefficiencies in the Monte-Carlo generation of the processes of interest, the initial state radiation parton showers are added with a backward evolution technique. Starting from the hard interaction at parton level, the probabilities of the incoming partons to have originated from partons with higher momentum fraction x at lower scales are successively calculated until again a low cutoff scale is reached. The emitted partons obtained by this procedure are showered as described above for the case of final state parton showers.

Since both the matrix element calculation of the hard process and parton shower evolution are combined in the event generation, the overlap between the two approaches has to be removed to avoid double-counting of particular processes. Such overlap can arise if e.g. the emission of a parton during the shower evolution leads to an extra jet in the event which is already accounted

for by the matrix element calculation of the hard process. There are several matching schemes such as the MLM [121] or CKKW [122] procedures which remove this overlap. Details on the CKKW procedure can be found in the precited reference. In the MLM approach, partons generated at matrix element level are required to fulfill criteria on the transverse momentum $p_T^{\text{part}} > p_T^{\text{min}}$, the pseudo-rapidity $|\eta_{\text{part}}| < \eta_{\text{max}}$, and their spatial separation $\Delta R_{ij} > R_{\text{min}}$. The partons are then associated to jets which are formed from the parton showers using a cone jet algorithm. The event is accepted if all partons at matrix element level can be associated in ΔR with a parton shower jet and vice versa. The procedure is applied to samples which are generated exclusively for a fixed number of light partons at matrix element level. An exception is the highest available parton multiplicity sample for which extra jets not matching matrix element level partons are allowed.

Hadronization and Decay

The parton shower development is followed by hadronization of the colour-charged final state partons to colour-neutral final state hadrons which may further decay to stable particles. The hadronization takes place in the non-perturbative regime at low energy scales where the strong coupling constant α_s is large. There are several phenomenological models for the formation of hadrons which are tuned to experimental data. In the Lund string model [123] the colour force between the partons is represented by a string. The energy inside the string increases linearly as the constituting partons are separated. Eventually the string breaks and a new $q\bar{q}$ pair is produced at the endpoints. The process is repeated iteratively and eventually leads to the formation of hadrons represented by string segments. In the cluster hadronization model [124] each gluon of the parton shower is split into a $q\bar{q}$ pair and clusters are built from colour-charge connected quark pairs. Clusters of high mass are split iteratively until hadrons are formed. The hadrons built during the hadronization process may be subject to further decays. The decay probabilities and lifetimes are known from experimental measurements.

Underlying Event

All additional activity in the event that is not associated to the hard process of interest is collectively termed underlying event. There are several mechanisms that contribute to the underlying event. The so-called beam remnants are the partons of the incoming protons that do not directly participate in the hard interaction of interest. They are, however, connected via the colour-charge to the rest of the event and thus constitute a part of the same hadronization system. In addition, semi-hard multiple parton interactions (MPI) may originate from remnant-remnant interactions. In high luminosity environments, like at the LHC, one also expects a significant number of pile-up interactions which arise from different proton-proton collisions occuring during the same proton bunch crossing.

Chapter 5

Measurement of the Inclusive Electron Cross-Section from Heavy-Flavour Decays with 2010 Data

This chapter presents the first physics analysis carried out within this thesis - the measurement of the inclusive cross-section for electron production from heavy-flavour decays using the first $\mathcal{L} = 1.3 \text{ pb}^{-1}$ of data collected with the ATLAS detector at 7 TeV center-of-mass energy in proton-proton collision during the 2010 data-taking campaign. The physics motivation for the measurement and an overview of the analysis is given in section 5.1. Section 5.2 reviews the theoretical tools for the calculation of heavy-flavour cross-sections. The data and simulated samples used in the analysis are listed in section 5.3 and the corresponding trigger selection is described in section 5.4. Sections 5.5 and 5.6 detail the selection criteria and signal extraction method used to determine the number of electrons from heavy-flavour decays. The efficiency to select these electrons and the cross-section unfolding technique are described in section 5.7. The final results of the measurement, including a full list of systematic uncertainties, and comparisons to theoretical predictions and complementary measurements are shown section 5.8. Section 5.9 concludes on the results of the analysis.

The author's main contribution to the analysis was the development and optimization of the signal extraction technique described in section 5.6, including the determination of the final signal yields for the cross-section measurement, and the evaluation of the associated systematic uncertainties. The author also contributed to the preparation of the data for the final analysis, and the optimization of the selection requirements. The efficiency measurements described in section 5.7, though using in part the same methodology developed in the signal extraction section, was performed by collaborators, as was the cross-section unfolding, and the study and comparison of the various theoretical predictions.

5.1 Introduction

With the commissioning of the LHC and the vast programme of physics measurements and searches foreseen during its lifetime, a good initial understanding of the fundamental objects and quantities measured by the ATLAS detector is indispensable. Furthermore, precise measurements of the known Standard Model physics processes allow to improve their theoretical description and form the basis of searches for new physics scenarios. The measurement of the inclusive cross-section for electron production from heavy-flavour decays presented in this chapter aims to serve both of the above purposes.

Being one of the first analyses performed at the LHC, an important objective of the measurement was to gain an initial understanding of the electron reconstruction performance in the low transverse momentum range, where the spectrum is dominated by non-isolated electrons from heavy-flavour decays in addition to large backgrounds from hadrons faking electrons and background electrons from photon conversions. This knowledge is an important step towards searches for new physics scenarios where the signal acceptance and sensitivity can be significantly improved by using low- p_T electrons in addition to those observed at higher transverse momenta. Examples of searches where low- p_T electrons have been used successfully in ATLAS to date is the recent discovery of a new particle in the search for the Standard Model Higgs boson [40, 125] and, in the context of this thesis, the searches for compressed supersymmetric scenarios presented in chapter 6.

Due to the abundance of semi-leptonic heavy-flavour decays at low lepton transverse momenta, the measurement of the inclusive electron cross-section in this regime can be used to evaluate the accuracy of theoretical predictions for bottom and charm production and to constrain their uncertainties. The results are thus valuable input for the validation and further development of theoretical calculations which may themselves be used to predict backgrounds in future new physics searches. The need for testing heavy-flavour QCD predictions at the LHC becomes particularly clear when looking at the previous experience at the Tevatron, where an excess in the production of bottom quarks was observed (e.g. [126, 127]) that could only later be resolved with improved measurements (e.g. [128]) and more precise theoretical calculations [129, 130].

The electrons of interest to the analysis presented in this chapter originate from bottom and charm quarks which are produced via the strong interactions. Figure 5.1 shows the production processes for heavy quarks at the leading order of perturbation theory as well as a selection of important next-to-leading order corrections. In the process of hadronization the bottom and charm quarks form B - and D -hadrons which then further decay to electrons with branching fractions of approximately 10 %. Possible decay modes are $B \rightarrow e$ (dominant at $p_T^{\text{electron}} \gtrsim 5$ GeV), $D \rightarrow e$ (dominant at $p_T^{\text{electron}} \lesssim 5$ GeV), and the secondary decay $B \rightarrow D \rightarrow e$ (strongly suppressed at $p_T^{\text{electron}} > 1$ GeV) [131]. The electrons produced in semi-leptonic heavy-flavour decays are inherently non-isolated due to the nearby activity from hadron decays. A description of the state-of-the-art theoretical predictions for the electron production processes described above is given in section 5.2.

The measurement of the inclusive cross-section for electrons from heavy-flavour decays is performed in the electron transverse momentum range $7 < p_T < 26$ GeV and within pseudorapidities $|\eta| < 2$, excluding $1.37 < |\eta| < 1.52$. The measurement is not extended to higher transverse momenta since the contribution of isolated electrons from $W/Z/\gamma^*$ processes, which have different reconstruction and identification efficiencies, becomes dominant in that region.

Figure 5.1: (a)-(d): Leading order heavy quark production processes. (e)-(j): Selection of next-to-leading order heavy quark production processes. The diagrams (h)-(j) are referred to as gluon radiation, flavour excitation, and gluon splitting and contribute large next-to-leading logarithmic contributions at high heavy quark transverse momenta.

The differential cross-section is determined with the following expression:

$$\frac{\Delta\sigma_i}{\Delta p_{T_i}} = \left(\frac{N_i^{Q \rightarrow e}}{\epsilon_{\text{trigger}_i} \cdot \int \mathcal{L} dt} - \sigma_{\text{accepted}_i}^{W/Z/\gamma^*} \right) \cdot \frac{C_{\text{migration}_i}}{\epsilon_{(\text{reco}+\text{ID})_i} \cdot \Gamma_{\text{bin}_i}}. \quad (5.1)$$

The index i labels the p_T bins of the measurement with width Γ_{bin_i} . The integrated luminosity of the measurement is denoted $\int \mathcal{L} dt$ and corresponds to the first 1.3 pb⁻¹ of data taken during the 2010 data-taking campaign, where the low luminosity conditions allowed to select non-isolated low- p_T electrons without the application of trigger prescale factors or further high-level trigger isolation requirements. The corresponding trigger efficiencies are denoted $\epsilon_{\text{trigger}_i}$ and are measured on data. $N_i^{Q \rightarrow e}$ corresponds to the number of electrons from heavy-flavour decays which is extracted from the data with an extended maximum likelihood fit. The reconstruction and identification efficiencies, denoted $\epsilon_{(\text{reco}+\text{ID})_i}$, and the bin-to-bin migration correction for the cross-section unfolding, $C_{\text{migration}_i}$, are determined from simulation. The uncertainties on the identification efficiencies are derived from a direct data measurement. Finally, the small contribution from $W/Z/\gamma^*$ processes to the measured cross-section is subtracted.

The determination of these inputs to the cross-section calculation is described in sections 5.3-5.7. The final results are presented in section 5.8.

5.2 Theoretical Predictions

Amongst the state-of-the-art tools for the theoretical prediction of heavy flavour production is the FONLL program [129, 130]. FONLL provides calculations at fixed next-to-leading order accuracy with next-to-leading log resummation. In FONLL, the calculation of the inclusive cross-section for heavy-flavour production with subsequent leptonic decays is decomposed according to

$$\frac{d\sigma^e}{dp_T} = \frac{d\sigma^Q}{d\hat{p}_T} \otimes D_{\text{NP}}^{Q \rightarrow H_Q} \otimes g_{\text{weak}}^{H_Q \rightarrow l}, \quad (5.2)$$

where the first term is the heavy quark production cross-section evaluated with perturbative QCD at NLO+NLL accuracy (see below), the second term is the non-perturbative fragmentation function modeling the formation of heavy flavour hadrons, and the third term describes the weak decay into the observable lepton.

In the perturbative part of the calculation two regimes have to be considered. For heavy quark transverse momenta $p_T \approx m_Q$ the cross-section is accurately specified by the fixed order NLO formalism. In the case of $p_T \gg m_Q$, however, it is necessary to resum to next-to-leading accuracy (NLL) large logarithms of the form $\alpha_s^n \log^n(p_T/m_Q)$ and $\alpha_s^{n-1} \log^n(p_T/m_Q)$ which arise from collinear terms associated to gluon radiation, flavour excitation, and gluon splitting processes shown in Fig. 5.1 (h)-(j). In FONLL the NLO+NLL calculation in the high- p_T regime is performed by convolving the parton level cross-section σ_i^{NLO} with a so-called perturbative fragmentation function $D_i^Q(z, \mu_F)$ according to (e.g. [132])

$$d\sigma^Q \sim \int dz d\hat{\sigma}_i^{\text{NLO}} \times D_i^Q(z, \mu_F), \quad (5.3)$$

Figure 5.2: *Left: Uncertainties of the FONLL theoretical prediction as a function of the true electron transverse momentum in bottom and charm production for $|\eta| < 2.0$ excluding $1.37 < |\eta| < 1.52$. Right: Ratio of cross-sections predicted by the NLO-only calculation of FONLL (see text for details) as well as a number of Monte-Carlo generators, to the full FONLL NLO+NLL prediction as a function of true electron transverse momentum. Taken from reference [134].*

where i is the parton produced in the interaction, z the momentum fraction of the heavy quark, and μ_F the factorization scale. Here the parton density functions of the incoming hadrons have been suppressed for simplicity. The idea behind equation 5.3 is that the quark production cross-section at high scales $\mu_F \sim p_T \gg m_Q$ can be factorized into a NLO calculation in the limit of massless partons and the perturbative fragmentation function which describes the formation of the heavy quark. The perturbative fragmentation function can be determined with QCD calculations as demonstrated in [133]. FONLL performs the cross-section calculation in both transverse momentum regimes by appropriately matching the fixed order formalism with the NLL corrections. The implementation of this matching procedure is described in [130].

The hadronization of the heavy quarks is modelled with the non-perturbative fragmentation functions documented in [135, 136], the parameters of which are determined from fits to e^+e^- collider data [137–139] using NLO+NLL theory calculations for the $e^+e^- \rightarrow \text{jets}$ processes. The final weak decays to the observable lepton are obtained from decay spectra measured at B -factories using hadron masses and branching ratios as tabulated by the Particle Data Group [2].

The dominant uncertainty of the FONLL prediction originates from the renormalization and factorization scales. The central values of the scales are chosen as $\mu_0 = \mu_{R,F} = \sqrt{p_T^2 + m_Q^2}$, and the uncertainties are determined from scale variations within $0.5 < \mu_{R,F}/\mu_0 < 2.0$ while keeping the relation $0.5 < \mu_R/\mu_F < 2.0$. Further uncertainties arise from the heavy quark masses which are set to $m_b = 4.7 \pm 0.25$ GeV and $m_c = 1.5 \pm 0.2$ GeV and have the largest impact at low transverse momenta. The PDF uncertainties are the ones provided with the CTEQ6.6 [140] set which is also used for the central values of the prediction. Systematic uncertainties related to the value of the strong coupling constant within $0.116 < \alpha_s < 1.120$ are estimated with the CTEQ6.6AS [141] PDF set and are found to be negligible. The uncertainties associated with the choice of scales, quark masses, and PDFs are added in quadrature to obtain the final uncertainty. Fig. 5.2 (left) shows the uncertainties of the FONLL prediction in the

measurement range which is largest at low electron transverse momentum ($\sim 38\%$) due to the p_T dependence of the scale and quark mass uncertainties. Uncertainties related to the non-perturbative fragmentation function are not taken into account. They are expected to be small, typically around a few percent, reaching about 5 % at low transverse momentum [142].

For comparison, Fig. 5.2 (right) shows the ratio of the NLO-only¹ prediction of FONLL with respect to the full FONLL prediction. It can be seen that there is no real sensitivity to the NLL resummation within the measurement range since both predictions agree within their uncertainties. However, the theoretical uncertainties at higher transverse momenta are reduced if the NLL resummation is included into the calculation.

5.3 Data and Simulated Samples

5.3.1 Simulated Samples

The simulated signal and background samples used in this analysis are listed in Tab. 5.1. The $b\bar{b}$ and $c\bar{c}$ signal samples are produced with filters on the transverse momentum of the leading electron at generator level ($p_T^{\text{el},1} > 3, 5, \text{ and } 7 \text{ GeV}$) to provide sufficient statistics over the full transverse momentum range of the analysis. An additional filter is applied to the $c\bar{c}$ samples to remove events with electrons originating from b -decays with intermediate D -hadron production. Additional samples modeling J/ψ and $W/Z/\gamma^*$ production processes with filters at generator level as indicated in Tab. 5.1 are used to estimate the contributions from isolated electrons. The dominant background processes from hadrons faking electrons and photon conversions are contained in the minimum bias and di-jet samples. The minimum bias sample contains the low- p_T production and semi-hard $2 \rightarrow 2$ QCD processes listed in [143]. The di-jet sample is dominated by QCD di-jets, but is generated inclusively with heavy flavour, prompt photon, and W/Z production processes. The minimum bias and di-jet samples are filtered at generator level requiring a minimum transverse energy threshold of the generated particles within a geometrical window in η and ϕ as given in Tab. 5.1. In the analysis the transition between the 3 GeV and 7 GeV filtered PYTHIA $b\bar{b}$ and $c\bar{c}$ sample is made at 14 GeV reconstructed electron p_T . The transition between the minimum bias and di-jet samples is made at 16 GeV reconstructed electron p_T for photon conversions, and 18 GeV for hadrons faking electrons. The continuity of the observables of the measurement at these thresholds has been validated.

The baseline samples used in the analysis are exclusively generated with PYTHIA 6.421 [143] at LO accuracy. For systematic studies and comparisons with the measurement further samples of electrons from heavy-flavour decays are generated with POWHEG-hvq 1.0 patch 4 [144, 145] at NLO accuracy, interfaced to either PYTHIA or HERWIG 6.510 [146]. In conjunction with HERWIG, JIMMY 4.31 [147] is used to model the underlying event. The POWHEG samples use PHOTOS 2.15 [148] to model final state QED radiation.

The PDF set used is MRST LO* [149] for the PYTHIA samples. The POWHEG+PYTHIA and POWHEG+JIMMY/HERWIG samples are used with CTEQ6.6 [140] at matrix element level and CTEQ6L1 [150] for parton showering and underlying event. All samples are generated with the ATLAS MC09 tune [151] and are passed through the full chain of ATLAS event simulation described in section 4.4.

¹NLO-only here refers to the calculation at matrix-element level. The default FONLL non-perturbative fragmentation function obtained from NLO+NLL fits of e^+e^- data is used.

Physics process	Generator	Filter	N^{gen}	$\sigma^{\text{gen}} \times \epsilon^{\text{gen}}$ [pb]
$b\bar{b}$	PYTHIA	$p_T^{\text{el},1} > 3 \text{ GeV}$	4M	17.2×10^6
$c\bar{c}$	PYTHIA	$p_T^{\text{el},1} > 3 \text{ GeV}$	4M	19.8×10^6
$b\bar{b}$	PYTHIA	$p_T^{\text{el},1} > 7 \text{ GeV}$	4M	1.79×10^6
$c\bar{c}$	PYTHIA	$p_T^{\text{el},1} > 7 \text{ GeV}$	4M	0.96×10^6
$b\bar{b}$	POWHEG+PYTHIA	$p_T^{\text{el},1} > 5 \text{ GeV}$	0.3M	2.02×10^6
$c\bar{c}$	POWHEG+PYTHIA	$p_T^{\text{el},1} > 5 \text{ GeV}$	0.3M	1.60×10^6
$b\bar{b}$	POWHEG+JIMMY/HERWIG	$p_T^{\text{el},1} > 5 \text{ GeV}$	0.3M	1.26×10^6
$c\bar{c}$	POWHEG+JIMMY/HERWIG	$p_T^{\text{el},1} > 5 \text{ GeV}$	0.3M	1.45×10^6
$W \rightarrow e\nu$	PYTHIA	-	7M	8.93×10^3
$Z \rightarrow ee$	PYTHIA	-	5M	0.856×10^3
$J/\Psi \rightarrow ee$	PYTHIA	$p_T^{\text{el},1} > 3 \text{ GeV}$	5M	39.16×10^6
Drell-Yan $\rightarrow ee$	PYTHIA	$1 < M_{ee} < 60 \text{ GeV}, p_T^{\text{el},2} > 3 \text{ GeV}$	1M	2.36×10^3
Drell-Yan $\rightarrow ee$	PYTHIA	$8 < M_{ee} < 15 \text{ GeV}, p_T^{\text{el},1} > 3 \text{ GeV}$	0.5M	2.21×10^3
Drell-Yan $\rightarrow ee$	PYTHIA	$15 < M_{ee} < 60 \text{ GeV}$	0.8M	1.25×10^3
Minimum bias	PYTHIA	$E_T^{\text{true}} > 6 \text{ GeV}$ within $\eta \times \phi = 0.12 \times 0.12$	40M	2.083×10^9
Di-jets	PYTHIA	$E_T^{\text{true}} > 17 \text{ GeV}$ within $\eta \times \phi = 0.12 \times 0.12$	40M	97.6×10^6

Table 5.1: Summary of simulated samples used in the analysis, indicating for each physics process, the Monte-Carlo generator, the filter criteria applied during generation, the number of events generated, and the resulting value of the LO (PYTHIA) or NLO (POWHEG) cross-section times filter efficiency. The minimum bias sample contains low- p_T and semi-hard QCD production processes. The di-jet sample is dominated by QCD di-jets, but is generated inclusively with heavy-flavour, prompt photon, and W/Z production processes. Details on the implementation of these processes can be found in [143]. The minimum bias and di-jet samples are filtered at generator level requiring a minimum transverse energy threshold within a geometrical window in η and ϕ .

Fig. 5.2 (right) shows the cross-section predictions of the heavy-flavour simulated samples listed above divided by the FONLL prediction as a function of electron transverse momentum. Significant differences in the cross-section normalizations are observed for the PYTHIA and POWHEG+HERWIG/JIMMY simulated samples. Since PYTHIA is a LO generator it is not expected to match the NLO predictions. Additional sources that may contribute to the discrepancies of both generator configurations are the underlying parton shower and hadronization models, the generator tunes, and the B - and D -hadron decay models. The effect of the latter has been studied in detail by comparing the cross-section predictions of the standard PYTHIA and HERWIG decay tables to the predictions of EVTGEN [152]. Fig. 5.3 (left) indicates that differences of up approximately 10 % are expected between PYTHIA and EVTGEN, whereas the HERWIG prediction is approximately 30-40 % lower than EVTGEN. Furthermore, the effect of the different generator tunes was studied by switching off the simulation of multiple parton interactions in the modeling of the underlying event. This modification, however, did not result into a significant change in the predicted cross-section. In addition, the effect of final state QED radiation (FSR) on the predicted cross-section was studied using a number of different methods for dressing electrons with photons. The results of this study are shown in Fig. 5.3 (right). The maximum variations observed are between 3 % and 5 % inside the measurement range.

Figure 5.3: Ratios of electron cross-sections from heavy-flavour decays as a function of generated electron p_T within $|\eta| < 2$, excluding $1.37 < |\eta| < 1.52$ using PYTHIA, POWHEG+PYTHIA, and POWHEG+HERWIG/JIMMY generated samples. Left: Cross-section ratios using the EVTGEN and nominal decay tables of the corresponding generators, where either B- and D-hadrons or only B-hadrons are re-decayed in EVTGEN. Right: Cross-sections ratios before and after QED final state radiation (FSR) using PHOTOS with the following configurations: Electron four-momentum correction with photons coming from the same vertex within $\Delta R < 0.1$ (GenVtx $\Delta R < 0.1$), with photons coming from the same vertex and being closest in ΔR if more than one electron is associated to the vertex (GenVtx closest), or with photons from any vertex within $\Delta R < 0.1$ ($\Delta R < 0.1$). Taken from reference [134].

5.3.2 Data Sample

The analysis uses the first ATLAS data collected at 7 TeV centre-of-mass energy in proton-proton collisions during the 2010 data-taking campaign. Fig. 5.4 (left) shows the cumulative online luminosity distribution versus day throughout the year 2010 as delivered by the LHC (green) and as recorded by ATLAS (yellow). The total delivered luminosity amounts to 48.1 pb^{-1} of which 45 pb^{-1} were recorded by the ATLAS experiment. The peak instantaneous luminosity reached a value of approximately $2 \times 10^{32} \text{ cm}^{-2} \text{ s}^{-1}$ during that time. Table 5.2 gives a summary of the 2010 data-taking campaign which is subdivided into a number of periods labelled alphabetically from A to I according to the detector and LHC conditions. During the course of the year the operation parameters evolved to continuously higher luminosities starting with one colliding bunch in the LHC to a maximum of 348 colliding bunches at 150 ns bunch spacing. The distribution of the peak number of interactions per beam crossing as a function of time is shown in Fig. 5.4 (right). The values are significantly lower than during the subsequent data taking period in 2011 (see section 6.2.3).

Given the low transverse momentum range of electrons used in the analysis and the evolution of the trigger thresholds with the increasing instantaneous luminosity provided by the LHC throughout the year, only the first 1.3 pb^{-1} of data are used in the analysis. These data were taken between March and August 2010 in 82 detector runs that were deemed acceptable for physics analysis by the ATLAS data quality groups. The uncertainty on the integrated luminosity amounts to 0.04 pb^{-1} and is determined as described in [153].

Period	Date Range	Run Range	Stable / Ready Lum. [pb ⁻¹]	Max. Inst. Lum. [10 ³⁰ cm ⁻² s ⁻¹]	$\mu_{\max}^{\text{avg}}$	Fill Range	Bunch Count	Bunch Δt [ns]
A	Mar-30 - Apr-19	152166 - 153200	4.3e-4 / 4.0e-4 (92%)	2.5e-3	1.6e-2	1003 - 1049	1	-
B	Apr-23 - May-17	153565 - 155160	8.9e-3 / 8.7e-3 (98%)	6.5e-2	1.4e-1	1057 - 1107	1 - 3	5000 - 44625
C	May-17 - Jun-05	155228 - 156682	8.4e-3 / 7.6e-3 (91%)	2.2e-1	1.8e-1	1108 - 1134	3 - 8	2500 - 5000
D	Jun-24 - Jul-19	158045 - 159224	3.2e-1 / 3.1e-1 (97%)	1.5	2.1	1178 - 1233	2 - 8	2500 - 44625
E	Jul-29 - Aug-18	160387 - 161948	1.1 / 1.1 (97%)	3.8	1.5	1251 - 1293	16	2350 - 2500
F	Aug-19 - Aug-30	162347 - 162882	2 / 1.9 (97%)	9.9	2	1295 - 1309	32 - 36	1000 - 1250
G	Sep-22 - Oct-07	165591 - 166383	9.4 / 9 (95%)	68	2.7	1364 - 1397	16 - 186	50 - 150
H	Oct-07 - Oct-18	166466 - 166964	9.1 / 8.6 (95%)	143	3.1	1399 - 1431	3 - 295	150 - 1250
I	Oct-24 - Oct-29	167575 - 167844	24 / 23 (96%)	198	3.6	1436 - 1453	295 - 348	50 - 150
Total	Mar-30 - Oct-29	152166 - 167844	46 / 44 (96 %)	198	3.6	1003 - 1453	1-348	50 - 44625

Table 5.2: *Summary of run periods with stable beams in proton-proton collisions at 7 TeV center-of-mass energy during the 2010 data-taking campaign. For each period important ATLAS detector and LHC parameters are given. “Stable luminosity” refers to the integrated luminosity over all stable beam luminosity blocks while “ready luminosity” corresponds to the subset of stable luminosity where all ATLAS detector systems were at their operational settings for physics data-taking. The difference to the recorded luminosity value in Fig. 5.4 is in part explained by the trigger deadtimes at L1. The quantity $\mu_{\max}^{\text{avg}}$ is the maximum of the luminosity block averaged values of the number of interactions per beam crossing in a given period. The column “bunch count” refers to the number of colliding bunches in the LHC. “Bunch Δt ” is the minimum time between colliding bunches. The values and ranges in this table are for illustrative purpose only and do not necessarily correspond to those during physics data-taking.*

Figure 5.4: *Left: Cumulative distribution of luminosity versus day delivered to (green), and recorded by ATLAS (yellow) during stable beams for proton-proton collisions at 7 TeV centre-of-mass energy in 2010. Right: Peak number of interaction per beam crossing versus day during the 2010 data-taking campaign. Further details in [153].*

5.4 Trigger Selection

The data sample was collected with four L1 electromagnetic triggers with progressively tighter transverse energy requirements as the instantaneous luminosity of the LHC increased. For each data-taking period the lowest unprescaled L1_EM item was chosen. No high-level electron triggers were used due to the shower shape and implicit isolation requirements encapsulated in the R_η and R_{had} variables which are part of all high level electron triggers (see Tab. 4.1) and significantly reduce the selection efficiency of non-isolated electrons from heavy-flavour decays.

Tab. 5.3 lists the triggers used in each data-taking period along with the applied transverse energy thresholds for offline reconstructed electron candidates in the analysis, and the absolute and fractional integrated luminosities. Approximately 1 % of the data sample was taken with the L1_EM2 item. The minimum transverse energy of reconstructed electron candidates recorded during this period was set to 7 GeV. In period D the L1_EM5 and L1_EM10 items were used to collect 9 % and 14 % of the total data respectively. The corresponding thresholds for offline

L1 Trigger	Period	E_T^{cluster} -threshold [GeV]	L^{int} [nb^{-1}]	L^{int} -fraction
L1_EM2	ABC	7	14.1 ± 0.4	1 %
L1_EM5	D	10	111.6 ± 3.5	9 %
L1_EM10	D	14	176.4 ± 5.6	14 %
L1_EM14	E	18	975.0 ± 31.2	76 %
Total	A-E	7-18	1277 ± 41	100 %

Table 5.3: *Summary of L1 trigger items per data taking period used in the analysis. The corresponding electron cluster transverse energy thresholds to be applied to reconstructed electron candidates (see section 5.5) as well as the integrated luminosities and luminosity fractions are shown.*

reconstructed electrons were set to 10 and 14 GeV. The bulk of the data (approximately 76 %) was collected in period E using the L1_EM14 item. The associated threshold for offline reconstructed electrons in this period was set to 18 GeV.

5.5 Electron Candidate Selection

The electron selection requirements are grouped into preselection and identification criteria. As part of the preselection, a primary vertex in the nominal beam spot region with at least three associated tracks is required. From these events electron candidates are reconstructed using the algorithm described in section 4.3.1. The electron candidates are required to fulfill the period-dependent transverse momentum requirements listed in Tab. 5.3 and to lie within the pseudorapidity coverage of the TRT, $|\eta| < 2.0$. Electrons candidates falling into the transition regions between the barrel and endcap calorimeters, $1.37 < |\eta| < 1.52$, are discarded. In addition to these criteria the tracks associated to electrons are required to have at least four hits in the silicon detectors and at least ten hits in the TRT. Electron candidates with clusters located close to problematic regions in the electromagnetic calorimeter or with associated tracks passing through inactive B-layer modules are rejected. In addition, the fraction of the raw cluster energy deposited in the strip layer of the electromagnetic calorimeter, f_1 , is required to be larger than 0.1.

For electron identification a selection of cuts from the nominal identification scheme described in section 4.3.1 is applied. The chosen criteria are listed in Tab. 5.4 and have been previously defined in Tab. 4.2. The selection reflects the need to optimize for non-isolated electrons from heavy-flavour decays and to use variables that are well-modelled in Monte-Carlo simulation.

Variable name	
$ \eta < 2.47$	•
$R_{\text{had}} \ \& \ R_{\text{had1}}$	
R_{η}	
$w_{\eta2}$	•
w_{stot}	•
E_{ratio}	•
n_{BL}	
n_{pixel}	•
n_{si}	•
d_0	•
d_0^{tight}	•
$\Delta\eta$	•
$\Delta\eta^{\text{tight}}$	
$\Delta\phi$	
E/p	
n_{TRT}	
f_{TR}	
Conversion Match	

Table 5.4: *Subset of criteria of the electron identification schemes described in section 4.3.1 used for electron identification in the analysis. The definition of the variables can be found in Tab. 4.3.*

Figure 5.5: *Distribution of electron candidates within $|\eta| < 2.0$, excluding $1.37 < |\eta| < 1.52$ as a function of cluster transverse energy in the range $7 < E_T < 100$ GeV (left) and pseudorapidity (right) in the range $7 < E_T < 26$ GeV for data and simulated samples. The data with $E_T < 18$ GeV in the left figure are scaled up from lower luminosities to 1.3 pb^{-1} to preserve the continuity of the spectrum. The W and Z/γ^* components are normalized to their NNLO cross-sections [154] and the remaining samples are scaled collectively from their respective LO cross-sections to match the data.*

The distribution of preselected electron candidates within $|\eta| < 2.0$, excluding $1.37 < |\eta| < 1.52$, is shown in Fig. 5.5 as a function of cluster transverse energy in the range $7 < E_T < 100$ GeV (left) and for pseudorapidity in the range $7 < E_T < 26$ GeV (right), for data and simulated samples where the data with $E_T < 18$ GeV has been scaled up from lower luminosities to 1.3 pb^{-1} to preserve the continuity of the spectrum. The simulated samples are the PYTHIA samples listed in Tab. 5.1, where the W and Z/γ^* components are normalized to their NNLO cross-sections [154] and the remaining samples are scaled collectively from their respective LO cross-sections to match the data. The discontinuities that are visible at the boundaries of 10, 20, and 30 GeV are due to the E_T -dependence of the electron identification cuts. The targeted E_T -range of the measurement presented in this thesis is $7 - 26$ GeV. The lower threshold is governed by the generator-level filter applied to the minimum bias sample used in the analysis (see Tab. 5.1). The measurement is not extended beyond 26 GeV since the fraction of isolated electrons from $W/Z/\gamma^*$ processes becomes significant in this regime and contaminates the sample of non-isolated electrons from heavy flavour decays which have different identification efficiencies. While the imperfect modeling of the identification efficiencies in simulation can be overcome for isolated electrons with data measurements using the W/Z resonances, the continuous spectrum of heavy flavour production makes an efficiency determination in the high- p_T range difficult.

It can be seen from Fig. 5.5 that the dominant background contributions in this region are from hadrons faking electron, (~ 70 % of observed electrons) and from photon conversions (~ 20 % of observed electrons). Consequently, the fraction of electrons from heavy-flavour decays amounts to ~ 10 %. The contribution of isolated electrons from W and Z/γ^* processes is small in the envisaged measurement range.

Tab. 5.5 shows the number of electron candidates observed in data after application of the preselection requirements in the nine bins of cluster transverse momentum in which the cross-section measurement is performed.

Period / L1 Trigger	Selected candidates	Selected candidates by E_T bin [GeV]								
		7-8	8-10	10-12	12-14	14-16	16-18	18-20	20-22	22-26
ABC / EM2	66 991	25387	24205	8102	3946	2014	1126	713	407	444
D / EM5	135 722	-	-	62904	30798	16116	8727	5506	3025	3637
D / EM10	64 398	-	-	-	-	24047	13539	8530	4819	5636
E / EM14	142 079	-	-	-	-	-	-	42904	25315	30766
Total	409 190	25387	24205	71006	34744	42177	23392	57653	33566	40483
L^{int} [nb ⁻¹]		14.1	14.1	125.7	125.7	302.1	302.1	1277	1277	1277

Table 5.5: *Numbers of electron candidates and corresponding integrated luminosities in bins of cluster transverse energy for each data-taking period [134].*

5.6 Signal Extraction

As described in the previous section the sample of selected electron candidates in the range foreseen for the cross-section measurement is dominated by hadrons faking electrons (hereafter denoted $h \rightarrow e$) and photon conversions (hereafter denoted $\gamma \rightarrow e$). The inability to obtain a sufficiently pure sample of signal electrons from heavy-flavour decays (hereafter denoted $Q \rightarrow e$) by applying additional selection criteria, necessitates the usage of a more sophisticated fit-based signal extraction technique. The implementation of this technique, the so-called Tiles Method, is described in section 5.6.1. Section 5.6.2 details several tests that have been performed to validate this method. The associated systematic uncertainties are discussed in section 5.6.3. The final results of the extraction are summarized in section 5.6.4.

5.6.1 Implementation of the Tiles Method

The fraction of $Q \rightarrow e$ electrons in the sample of electron candidates amounts to $\sim 10\%$. There are additional powerful discriminating variables that can be used to increase this fraction, most notably the fraction of high threshold transition radiation hits in the TRT out of all TRT hits measured on the electron track, f_{TR} , the number of hits in the pixel B-layer, n_{BL} , and E/p , the ratio of the electron cluster energy and the reconstructed track momentum. The distributions of these variables for electron candidates within the range $7 < E_T < 26$ GeV are shown in Fig. 5.6. The fraction f_{TR} provides discrimination against the $h \rightarrow e$ component for which no transition radiation is expected. In E/p a wider distribution is expected for the $h \rightarrow e$ and $\gamma \rightarrow e$ components, and $h \rightarrow e$ reaches its maximum at values smaller than one. The n_{BL} distribution allows to separate the $\gamma \rightarrow e$ component for a subset of candidates where a pixel B-layer is expected.

The application of appropriate selection requirements on the f_{TR} , E/p , and n_{BL} variables allows the fraction of $Q \rightarrow e$ electrons to be further increased in the candidate sample to approximately 50 %, which is clearly insufficient for the envisaged measurement and leaves no means (i.e. discriminating variables) to further estimate the remaining backgrounds. For this reason a multi-dimensional fit-based approach was adopted for the extraction of the $Q \rightarrow e$ yield.

In a first attempt a standard maximum likelihood fit was employed taking the f_{TR} and n_{BL} probability distribution functions (pdfs) for each component from simulation. However, given the dominance of the $h \rightarrow e$ background, studies showed that any small difference between data

Figure 5.6: *Distributions of the discriminating variables used in the signal extraction. The simulated samples are normalized collectively from their LO cross-sections to data. The fraction of high threshold transition radiation hits f_{TR} (top left), the number of hits in the pixel B-layer n_{BL} (top right), and the ratio of the electron cluster energy and the reconstructed track momentum E/p (bottom) are shown.*

and simulation in the pdfs for this component would result in large biases for the smaller $Q \rightarrow e$ signal component. Additionally, attempts to obtain a pure and unbiased $h \rightarrow e$ sample from the data by inverting preselection or identification criteria were spoiled by correlations between these variables and f_{TR} .

The solution to these difficulties was the introduction of the so-called “Tiles” method, a novel maximum-likelihood based approach, that was originally developed by the author and collaborators for background estimation in SUSY searches and is documented in the author’s diploma thesis [1] as well as reference [155]. The advantage of the Tiles Method is that under certain conditions it allows shape and normalization of one unknown component (in this case $h \rightarrow e$) to remain as free parameters of the fit. The implementation of the method in the context of this analysis is explained in the following.

The data is sliced into three-dimensional bins (or “tiles”) in the space spanned by the discriminating variables f_{TR} , n_{BL} , and E/p . The number of electrons $N(i, j, k)$ in a given bin identified with indices i, j, k is given by

$$N(i, j, k) = N^{Q \rightarrow e} p^{Q \rightarrow e}(i, j, k) + N^{\gamma \rightarrow e} p^{\gamma \rightarrow e}(i, j, k) + N^{h \rightarrow e} p^{h \rightarrow e}(i, j, k), \quad (5.4)$$

where the three N^α are the total number of electrons of each component α and the $p^\alpha(i, j, k)$ are the corresponding three dimensional pdfs for component α . The extended log-likelihood expression is then given by

$$-\ln L(N^{Q \rightarrow e}, N^{\gamma \rightarrow e}, N^{h \rightarrow e}) = \sum_{i,j,k} N(i, j, k) - N_{\text{obs}}(i, j, k) \ln N(i, j, k), \quad (5.5)$$

where $N_{\text{obs}}(i, j, k)$ is the number of observed candidates in bin i, j, k . The assumption is made that the unknown three-dimensional $h \rightarrow e$ pdf can be factorised as the product of the three one-dimensional f_{TR} , n_{BL} and E/p templates according to

$$p^{h \rightarrow e}(i, j, k) = p_{f_{\text{TR}}}^{h \rightarrow e}(i) \cdot p_{n_{\text{BL}}}^{h \rightarrow e}(j) \cdot p_{E/p}^{h \rightarrow e}(k), \quad (5.6)$$

where $p_{f_{\text{TR}}}^{h \rightarrow e}$, $p_{n_{\text{BL}}}^{h \rightarrow e}$, and $p_{E/p}^{h \rightarrow e}$ are the one-dimensional pdfs and are additional unknowns in the fit. Assuming that n_i , n_j , and n_k , are the number of bins in each observable the number of degrees of freedom associated to the $h \rightarrow e$ pdf is reduced from $n_i \times n_j \times n_k - 1$ to $(n_i - 1) + (n_j - 1) + (n_k - 1) = n_i + n_j + n_k - 3$ when going from the three-dimensional to the one-dimensional configuration.

In this particular implementation, $2(f_{\text{TR}}) \cdot 2(n_{\text{BL}}) \cdot 3(E/p) = 12$ bins are chosen in the discriminating variables. The bin boundary in f_{TR} is set to values between 0.05 and 0.15 in an E_{T} - and η - dependent way in order to optimize the sensitivity. The two bins in the number of pixel B-layer hits are defined as $n_{\text{BL}} = 0$ and $n_{\text{BL}} \geq 1$. In E/p the binning $E/p < 0.8$, $0.8 < E/p < 2.0$, and $E/p > 2.0$ is adopted.

With the twelve observations in the bins of the discriminating variables and the pdfs $p^{Q \rightarrow e}$ and $p^{\gamma \rightarrow e}$ taken from simulation (see Tab. 5.1), there are seven unknown parameters in the likelihood, namely the normalizations $N^{Q \rightarrow e}$, $N^{\gamma \rightarrow e}$, $N^{h \rightarrow e}$, and four shape parameters for the $h \rightarrow e$ pdfs given the above binning. These parameters are determined by maximizing equation 5.5 using the ROOT [96] implementation of MINUIT [156].

The extraction is performed independently in nine bins of electron E_{T} (7-8 GeV, 8-10 GeV, 10-12 GeV, 12-14 GeV, 14-16 GeV, 16-18 GeV, 18-20 GeV, 20-22 GeV, 22-26 GeV) and four bins of the absolute value of the pseudorapidity of the electron track $|\eta|$ (0.0-0.625, 0.625-1.07, 1.07-1.37, 1.52-2.0). The separation into pseudorapidity bins is motivated by the different f_{TR} response in these regions as shown in Fig. 3.7 of chapter 3.

5.6.2 Validation of the Tiles Method

A number of tests have been performed with simulated samples in order to validate and optimize the baseline configuration of the Tiles Method used for the signal extraction. The results of these tests are shown in this section for a selection of representative $(E_{\text{T}}, |\eta|)$ bins of the measurement. The full E_{T} - and η -dependent evaluation of the systematic uncertainties is described in section 5.6.3.

Correlation Bias

As mentioned in the previous section, the only potential bias intrinsic to the Tiles Method arises from possible correlations between the discriminating variables in the hadron component, which may spoil the assumption of equation 5.6 that the unknown three-dimensional $h \rightarrow e$ pdf can be factorized as the product of the three one-dimensional pdfs in f_{TR} , n_{BL} and E/p .

The effect of neglecting any correlations in the $h \rightarrow e$ component is evaluated by comparing the results of two closure tests. In the first test (“standard setup”) 5000 pseudo-experiments were performed, in each of which the number of electron candidates of each type in each bin of the $(f_{\text{TR}}, n_{\text{BL}}, E/p)$ distribution was randomly sampled from a Poisson distribution with mean equal to the expectation from simulation in 1.3 pb^{-1} of data. A second test (“uncorrelated setup”) was performed in an identical way, except that the correlations in the $h \rightarrow e$ component of the pseudo-data were removed via the factorization of the three-dimensional $(f_{\text{TR}}, n_{\text{BL}}, E/p)$ pdf into the product of the three one-dimensional pdfs of each variable.

The results of both tests are summarized in Fig. 5.7 for one particular bin in E_T and $|\eta|$. The absolute (top left), and relative residuals (top right) for the $Q \rightarrow e$ components are shown. It

Figure 5.7: Absolute (top left) and relative (top right) residuals as well as signed pull distributions (bottom left) and corresponding χ^2 -values (bottom right) obtained by running the Tiles extraction on 5000 pseudo-experiments with and without correlations in the $h \rightarrow e$ hadron component. This particular example shows the extraction in bin $7 \text{ GeV} < E_T < 8 \text{ GeV}$, $1.52 < |\eta| < 2.0$.

can be seen that the uncorrelated $h \rightarrow e$ configuration yields exact closure. Consequently, the small deviation from the true values observed in the standard setup ($\sim 6\%$ in this particular example) may be attributed to correlations that exist between the discriminating variables in the $h \rightarrow e$ component.

Fig. 5.7 (bottom left) shows the signed pull distributions, defined as the difference of extracted and true values of $N^{Q \rightarrow e}$ divided by the statistical uncertainty of the extracted event yield. It can be observed that the signed pull distributions follow a standard Gaussian with unit variance which confirms the accuracy of the error estimates obtained from the minimization.

The presence of correlations in the $h \rightarrow e$ component is also expected to affect the goodness-of-fit, which is measured in terms of

$$\chi^2 = \sum_{i,j,k} \frac{(N_{\text{obs}}(i,j,k) - N(i,j,k))^2}{N(i,j,k)} \quad (5.7)$$

where $N_{\text{obs}}(i,j,k)$ is the observed number of events in bin i,j,k of the three-dimensional ($f_{\text{TR}}, n_{\text{BL}}, E/p$) pdf and $N(i,j,k)$ is the expected number of events in this bin as returned by the fit. Figure 5.7 (bottom right) shows the distribution of χ^2 -values obtained from pseudo-experiments with and without correlations in the $h \rightarrow e$ component. For illustrative purposes an analytic χ^2 - distribution with $n_{\text{dof}} = n_{\text{obs}} - n_{\text{unknowns}} = 12 - 7 = 5$ degrees of freedom is overlayed, where n_{obs} is the number of observations and n_{unknowns} is the number of unknowns in the fit (see previous section). While excellent agreement is again demonstrated for the uncorrelated case the goodness-of-fit is expectedly deteriorated in the presence of correlations in the $h \rightarrow e$ component.

Linearity Test

The bias on the extracted number of signal events due to correlations in the $h \rightarrow e$ component may depend on the relative fractions of signal and backgrounds in data. This effect was studied by performing a so-called linearity test in which the signal extraction was performed on pseudo datasets constructed with different fractions

$$f_{\text{true}}^\alpha = \frac{N^\alpha}{N^{Q \rightarrow e} + N^{\gamma \rightarrow e} + N^{h \rightarrow e}} \quad (5.8)$$

of the signal and background components ($\alpha = Q \rightarrow e, \gamma \rightarrow e, h \rightarrow e$). The results of the test are shown in Fig. 5.8 for a particular E_{T} and $|\eta|$ bin. For a fixed fraction of either $h \rightarrow e$ (top) or $\gamma \rightarrow e$ (bottom), $f_{\text{true}}^{Q \rightarrow e}$ was varied in the pseudo-data generation. The extracted fractions of signal electrons $f_{\text{extracted}}^{Q \rightarrow e}$ and the corresponding relative biases as a function of the true fraction of signal electrons $f_{\text{true}}^{h \rightarrow e}$ are shown in the left and right panels of Fig. 5.8 respectively. In the top panels the $f_{\text{true}}^{h \rightarrow e}$ component was fixed to 50% (squares) and 80 % (triangles). In the bottom panels the $f_{\text{true}}^{\gamma \rightarrow e}$ component was fixed to 10% (squares) and 40% (triangles). The chosen values are conservative estimates based on simulation. One can see in the left panels that the departure from linearity (dotted line) is in general small. The right panels show that the corresponding bias is inversely proportional to the fraction of signal electrons and depends to a smaller extent on the fraction of the $\gamma \rightarrow e$ and $h \rightarrow e$ components.

Figure 5.8: Mean extracted fraction of signal electrons $f_{\text{extracted}}^{Q \rightarrow e}$ (left) and the corresponding relative bias (right) as a function of the true fraction $f_{\text{true}}^{Q \rightarrow e}$ obtained from 2000 pseudo-experiments for two constant values of $f_{\text{true}}^{h \rightarrow e}$ and $f_{\text{true}}^{\gamma \rightarrow e}$. This particular example shows the extraction in bin $10 < E_T < 12$ GeV, $1.52 < |\eta| < 2.0$.

Stability Test

The number of signal electrons $N^{Q \rightarrow e}$ extracted by the Tiles Method should be invariant with respect to the exact binning in the discriminating variables. However, systematic effects such as the correlations in the $h \rightarrow e$ pdf can break this invariance. Figure 5.9 shows the relative bias of the extraction and the associated statistical uncertainty as a function of the position of the inner bin boundaries of the continuous discriminating variables f_{TR} (two bins thus one inner bin boundary) and E/p (three bins thus two inner bin boundaries). Whereas some modulations within the statistical uncertainty of the extraction can be observed in the presence of correlations in the $h \rightarrow e$ component (left), the results are stable against variations of the inner bin boundaries if the correlations are removed (right). For variations of the second inner bin boundary of E/p correlations do not seem to have any notable impact. Based on the information given in Fig. 5.9, the position of the bin boundaries can be optimized such that the expected bias due to correlations in the $h \rightarrow e$ component as well as the statistical uncertainties remain small.

Figure 5.9: *Relative bias on the extracted number of signal electron candidates $N^{Q \rightarrow e}$ as a function of the position of the inner bin boundaries of f_{TR} (top) and E/p (middle and bottom) based on simulation. On the left side results are shown in the presence of correlations whereas on the right side correlations are removed. This particular example shows the extraction in bin $8 \text{ GeV} < E_T < 10 \text{ GeV}$, $0 < |\eta| < 0.625$.*

Figure 5.10: Comparison of the extracted number of signal electrons $N^{Q \rightarrow e}$ from the baseline Tiles Method with the results from the “ f_1 method” and the “ n_{TR} method”. For the former only the statistical uncertainty is shown. For the latter the full statistical and systematic uncertainties are included.

Comparison with Other Extraction Methods

The performance of the baseline signal extraction in the implementation described above has been compared to two alternative techniques referred to as the f_1 -method and the n_{TR} -method. The f_1 -method is based on the same Tiles formalism described in the previous section but interchanges the role of the f_1 and f_{TR} variables such that f_{TR} is used as a preselection requirement and f_1 as one of the discriminating variables in the fit. The n_{TR} -method is an independent approach that has no reliance on E/p and is based on the assumption that the number of high threshold TRT hits, n_{TR} , can be described by a Poisson distribution. The alternative methods are described in detail in [134]. Fig. 5.10 shows a comparison of the extraction results of the two methods and their corresponding uncertainties. It can be seen that all methods are compatible within their uncertainties. The baseline Tiles extraction provides the best experimental sensitivity².

5.6.3 Evaluation of Systematic Uncertainties

The signal extraction method is subject to various sources of systematic uncertainties, such as correlations in the $h \rightarrow e$ components, discrepancies between data and simulation in the f_{TR} , E/p , and n_{BL} pdfs of the $Q \rightarrow e$ and $\gamma \rightarrow e$ component, the determination of the electron energy scale, and the finite statistics of the simulated samples. The evaluation of these uncertainties is described in the following. The final uncertainties as a function of electron E_T as well as their combination is given at the end of this section.

²From Fig. 5.10 it may appear that the f_1 -extraction has the lowest uncertainties. However, the error bars represent only the statistical uncertainties. As will be shown in section 5.6.3 the f_1 -method suffers from larger systematic uncertainties than the baseline Tiles extraction due to correlations in the $h \rightarrow e$ component.

Correlations in the $h \rightarrow e$ component

The bias due to correlations in the $h \rightarrow e$ component in a given $(E_T, |\eta|)$ bin is governed by the nature of the correlations as well as the relative fractions of the $Q \rightarrow e$, $\gamma \rightarrow e$, and $h \rightarrow e$ components as demonstrated in section 5.6.2. The systematic uncertainty associated to the correlations is evaluated with 1000 pseudo-experiments per $(E_T, |\eta|)$ bin using the same procedure as described in section 5.6.2. In the construction of the pseudo-data, $f^{Q \rightarrow e}$, which is approximately constant as a function of E_T for a given bin in $|\eta|$, was conservatively taken as the lowest observed value in simulation across the full E_T -range of the measurement. The templates to model the $h \rightarrow e$ component in the pseudodata were taken from the filtered minimum bias ($7 \text{ GeV} < E_T < 18 \text{ GeV}$) and di-jet samples ($18 \text{ GeV} < E_T < 26 \text{ GeV}$) listed in Tab. 5.1. The impact of the limited statistics in these templates was evaluated with a so-called bootstrap technique [157]. From the original Monte-Carlo samples 250 “bootstrapped” templates were produced using the resampling with replacement method³. From each bootstrapped template a set of 1000 pseudo-experiments was generated and the corresponding correlation bias was evaluated. From the width of the distribution of these 250 measured values the statistical uncertainty on the estimated correlation bias was derived.

The results of this procedure are shown in Fig. 5.11 (blue markers) as a function of E_T for each of the four considered $|\eta|$ bins individually (top and middle) and their combination taking into account bin-to-bin correlations (bottom left). In the central pseudorapidity region the correlation bias is small and measured to good accuracy due to abundant Monte-Carlo statistics. At higher values of pseudorapidity larger fluctuations can be observed and a significant impact from the limited sample statistics is visible. In the combination, however, the higher pseudorapidity bins have a smaller weight since the bulk part of the $Q \rightarrow e$ component is located in the central region. Within the available statistics no strong E_T -dependence of the correlation bias inside the full $|\eta|$ -range of the measurement is observed. It can be seen clearly, however, that the uncertainty on the correlation bias estimate due to the limited Monte-Carlo sample statistics increases significantly with E_T between $7 \text{ GeV} < E_T < 18 \text{ GeV}$, and again after the transition to the di-jet sample between $18 \text{ GeV} < E_T < 26 \text{ GeV}$.

For the determination of the final uncertainty as a function of E_T two methods are considered. The first method (“method 1”) uses the maximum of the absolute central value of the bias estimate and its statistical uncertainty, $\max(|\langle \text{bias} \rangle|, \sigma_{\text{bias}})$, the second method (“method 2”) uses only the absolute central value, $|\langle \text{bias} \rangle|$. The uncertainties for both methods as a function of E_T are listed in Tab. 5.6. Given the large statistical uncertainties, and the fact that the bias estimation is purely Monte-Carlo driven, the first and more conservative method is chosen for the final uncertainty calculation. Since there is no obvious dependence on E_T within the uncertainties an averaging procedure as described at the end of section 5.6.3 is performed leading to a final uncertainty of 7.3 % across the full E_T -range.

Also shown in Fig. 5.11 (red markers) is the correlation bias for the previously introduced alternative construction of the Tiles Method which uses f_1 in place of E/p as discriminating variable in the fit. From the larger biases seen in this implementation (11.3 % using the procedure described above) it was concluded that the choice of E/p as discriminating variable is preferred.

³In the resampling with replacement method the events drawn from the original sample can appear more than once.

Figure 5.11: Relative bias due to correlations in the hadron component $h \rightarrow e$ as a function of E_T in bins of $|\eta|$ (top and middle) and for the combination of $|\eta|$ bins (bottom). The central values are obtained from 1000 pseudo-experiments and the statistical uncertainties are obtained with a bootstrapping technique as explained in text. Filtered minimum bias ($7 \text{ GeV} < E_T < 18 \text{ GeV}$) and di-jet samples ($18 \text{ GeV} < E_T < 26 \text{ GeV}$) are used for the construction of the hadron component in the pseudo-data. The blue markers show the baseline extraction method using E/p as one of the discriminating variables. The red markers show an alternative implementation using f_1 in place of E/p .

Discrepancies between Data and Simulated Samples

As described in section 5.6.1 the extraction method relies on $Q \rightarrow e$ and $\gamma \rightarrow e$ pdfs from simulated samples. Discrepancies between data and simulation in these templates may thus bias the extraction results. The systematic uncertainty associated to such discrepancies was evaluated individually for each discriminating variable in each $(E_T, |\eta|)$ bin as follows:

- f_{TR} : A high purity (approximately 95 %) sample of photon conversions was constructed as a subset of the sample of electron candidates by imposing $n_{BL} = 0$ and $E/p > 0.8$. The corresponding f_{TR} distribution over the full E_T - and η -range of the measurement is shown in Fig. 5.12 (left), where a discrepancy between data and simulation is visible. The associated systematic uncertainty was evaluated by modeling the observed difference as a shift in f_{TR} and applying it collectively to the $Q \rightarrow e$ and $\gamma \rightarrow e$ templates. The shift was applied in an E_T and $|\eta|$ dependent way, based on the observed difference between the means of the data and simulated distributions in that given bin. The range of values is -0.01 to 0.04. An example is shown in Fig. 5.13 for the transverse energy bin $8 \text{ GeV} < E_T < 10 \text{ GeV}$ and all four $|\eta|$ bins of the measurement, where a clear $|\eta|$ -dependence is visible. After application of the observed shift the signal extraction is repeated and the resulting change in $N^{Q \rightarrow e}$ is taken as the magnitude of the systematic uncertainty.
- E/p : Similarly as in the case of f_{TR} , a high purity (approximately 95 %) sample of conversions was constructed as a subset of the sample of electron candidates by imposing $n_{BL} = 0$ and $f_{\text{TR}} > 0.1$. The corresponding E/p distribution over the full E_T - and η -range of the measurement is shown in Fig. 5.12 (right). A discrepancy between data and simulation is visible which manifests itself in a shift of the distribution. The effect of this discrepancy is reduced in the extraction by placing the E/p bin boundaries in the pdfs away from the peak at $E/p \approx 1$. The systematic uncertainty is evaluated by conservatively applying a shift of 0.1 collectively to the $Q \rightarrow e$ and $\gamma \rightarrow e$ templates in all E_T and η bins of the measurement and repeating the signal extraction. The uniform shift of 0.1 covers well the observed differences between data and simulation in E_T and $|\eta|$, as shown for example in Fig. 5.14 for the transverse energy bin $10 \text{ GeV} < E_T < 12 \text{ GeV}$. The $|\eta|$ -dependence of the observed shifts is less pronounced as in the case of f_{TR} .

Figure 5.12: Comparison between data and simulation for f_{TR} (left) and E/p (right) for a pure sample of conversions, selected from the data using the criteria of section 5.5 with the additional requirements of $n_{BL} = 0$ together with $E/p > 0.8$ (left) and $\text{TRfrac} < 0.1$ for (right).

- n_{BL} : The correct simulation of the number of hits in the pixel B-layer is highly dependent on the accurate modeling of the beam pipe and ID material, where photon conversions are induced. Given the excellent understanding of the material in the ID based on the analysis of charged tracks and of K_s^0 decays in minimum-bias events [158], the following procedure is applied: The content of the two bins in n_{BL} for the $\gamma \rightarrow e$ pdf is coherently varied by $\pm 5\%$ for $n_{BL} = 0$ and $\mp 1\%$ for $n_{BL} > 0$. Again the signal extraction is repeated and the systematic uncertainty is determined as the maximum difference with respect to the nominal extracted event yield. The systematic uncertainty associated to the n_{BL} pdf of the $Q \rightarrow e$ component is estimated to be negligible since tracks going through non-functional B -layer modules are not considered in the analysis.

Figure 5.13: Shape comparison between data and simulation in f_{TR} for a high-purity sample of photon conversions in the transverse energy bin $8 \text{ GeV} < E_T < 10 \text{ GeV}$ and the four $|\eta|$ -bins of the measurement. The shift between the data and simulation is quantified with “ ΔMeans ”, the difference between the means of the distributions. An $|\eta|$ -dependence of ΔMeans can be observed.

The systematic uncertainties resulting from the shape variations described above are combined in $|\eta|$ taking into account possible bin-to-bin correlations. The results are summarized in Tab. 5.6 as a function of E_T . For the final result the uncertainties are averaged in η as further described at the end of this section.

To study the possible effects of pile-up interactions, which are not modelled in the simulation, the f_{TR} , n_{BL} and E/p distributions from the data collected in period E were remade with the requirement of exactly one vertex with more than two associated tracks. Within the small statistical errors, no difference in the shapes of the templates could be seen when compared to the default event selection. Therefore it was concluded that the small level of pile-up present in the data has a negligible effects on the signal extraction results.

Figure 5.14: Shape comparison between data and simulation in E/p for a high-purity sample of photon conversions in the transverse energy bin $10 \text{ GeV} < E_T < 12 \text{ GeV}$ and the four $|\eta|$ -bins of the measurement. The shift between the data and simulation is quantified with “ $\Delta \text{ Means}$ ”, the difference between the means of the distributions.

Electron Energy Scale

The electron energy scale in data is corrected according to the recommendations of the ATLAS electron / photon combined performance group implemented in the **EnergyRescaler** package [159]. The correction factors provided by this package are determined from measurements of the Z and J/ψ resonances, and in-situ studies comparing the ID track momenta with the measured calorimeter energies. The associated uncertainty is obtained by scaling the electron energy up and down by 1 % in the barrel and 3 % in the endcap regions of the electromagnetic calorimeter and repeating the signal extraction. The resulting changes in the event yields are shown in Tab. A.1 of Appendix A.1 in bins of E_T summed over $|\eta|$. The final uncertainties, taken as the absolute value of the maximum bias per E_T -bin, are shown in Fig. 5.6.

Statistics of Simulated Samples

For the determination of the uncertainties resulting from the finite Monte-Carlo statistics of the $Q \rightarrow e$ and $\gamma \rightarrow e$ pdfs a similar bootstrap technique to that previously described in this section is used. From the original Monte-Carlo samples 250 bootstrapped templates were produced using the resampling with replacement method. The signal extraction was repeated in each E_T and $|\eta|$ bin using these bootstrapped templates. The uncertainty was derived from the width of the resulting distribution of extracted $N^{Q \rightarrow e}$ values. Tab. A.2 in Appendix A.1 summarizes the full results as a function of E_T and $|\eta|$. The final uncertainties in bins of E_T are shown in Tab. 5.6. These values have a genuine E_T -dependence and are therefore not further averaged.

Summary of Systematic Uncertainties

Table 5.6 summarizes the systematic uncertainties discussed above as a function of E_T . Except for the uncertainty related to the limited Monte-Carlo statistics in the pdfs, all systematic uncertainties are averaged over E_T . The average is calculated by weighting the bias in each E_T -bin by the statistical uncertainty of the signal extraction in that bin. This procedure ensures that less weight is given to measurements that are associated to larger statistical uncertainties.

For the systematic uncertainty related to correlations in the $h \rightarrow e$ component the averaging in E_T is well motivated, since no E_T -dependence of the estimated bias within its statistical uncertainty is observed as demonstrated in the first part of this section. Table 5.6 shows the E_T -averaged bias values for both of the previously introduced methods. The more conservative “method 1” is chosen for the final uncertainty calculation.

The smaller systematic uncertainties related to the electron energy scale and the mismodeling of the f_{TR} , E/p , n_{BL} pdfs, are also averaged in E_T . The small E_T -dependence that may be present in some of these systematic uncertainties is not taken into account. This assumption has negligible impact on the final uncertainty of the measurement, which is dominated by the correlation uncertainty described above and uncertainties related to the determination of the reconstruction and identification efficiencies (see section 5.7).

The final uncertainties that are propagated to the cross-section measurement are marked in bold font in Tab. 5.6. Some of these uncertainties are correlated to those associated to the measurement of the identification efficiency described in section 5.7.3. These correlations are taken into account, when all uncertainties of the measurement are combined in section 5.8.2. The combined uncertainty associated to the signal extraction amounts to 9.1-9.4 %.

E_T bin [GeV]	7-8	8-10	10-12	12-14	14-16	16-18	18-20	20-22	22-26	Combined
f_{TR} -mismodeling	2.0	3.4	2.1	2.2	3.2	2.9	0.9	1.7	0.9	2.3
E/p -mismodeling	2.1	0.2	1.0	1.8	4.3	5.7	4.0	4.5	3.6	3.0
n_{BL} -mismodeling	2.2	2.1	1.5	1.2	1.0	1.1	0.6	0.8	0.7	1.3
Electron energy scale	4.3	4.8	3.1	3.3	4.8	3.5	2.2	3.5	3.4	3.5
Pdf statistics	0.9	0.8	1.4	1.8	2.5	1.4	2.0	2.5	2.5	0.8-2.5
$h \rightarrow e$ correlations, method 1	2.7	2.9	5.3	8.0	11.6	17.3	6.0	5.5	4.2	7.3
$h \rightarrow e$ correlations, method 2	1.1	0.6	0.3	2.4	8.2	2.7	6.0	5.5	0.5	2.9

Table 5.6: Summary of systematic uncertainties on the extracted number of signal electrons. The uncertainties are shown in percent for each E_T -bin of the measurement and, where applicable, for their combination. For the evaluation of the correlation bias two methods are considered as described further in the text. The numbers in bold font are the final uncertainties propagated to the measurement.

5.6.4 Extraction Results

The results of the signal extraction are shown in Tab. 5.7 (top) in bins of E_T and $|\eta|$ including statistical uncertainties. The bottom part of Tab. 5.7 shows the final results of the extraction together with the associated statistical and systematic uncertainties as a function of E_T after combination in $|\eta|$. For information the extracted $N^{h \rightarrow e}$ and $N^{\gamma \rightarrow e}$ yields including their statistical uncertainty are shown as well. The event yields are visualized and compared to the expectation from simulation in Fig. 5.15 as a function of E_T (top four panels) and $|\eta|$ (bottom

E_T -range [GeV]	$ \eta < 0.625$	$0.625 < \eta < 1.07$	$1.07 < \eta < 1.37$	$1.52 < \eta < 2.0$
7 – 8	1167 ± 48	831 ± 79	314 ± 67	622 ± 55
8 – 10	1100 ± 49	809 ± 64	512 ± 47	562 ± 55
10 – 12	3934 ± 103	2988 ± 1304	1806 ± 157	1610 ± 214
12 – 14	1891 ± 72	1374 ± 92	784 ± 80	864 ± 112
14 – 16	2134 ± 72	1686 ± 88	1079 ± 71	1102 ± 93
16 – 18	1195 ± 52	855 ± 61	541 ± 55	582 ± 63
18 – 20	2850 ± 91	2278 ± 1103	1162 ± 96	1201 ± 161
20 – 22	1708 ± 66	1418 ± 84	1020 ± 69	611 ± 86
22 – 26	1959 ± 70	1618 ± 90	958 ± 82	668 ± 122

E_T -range [GeV]	7-8	8-10	10-12	12-14	14-16	16-18	18-20	20-22	22-26
$N^{Q \rightarrow e}$	2934	2982	10337	4912	6000	3174	7490	4757	5203
$\sigma_{\text{stat.}}$	± 127	± 108	± 313	± 180	± 163	± 116	± 236	± 154	± 186
$\sigma_{\text{syst.}}$	± 267	± 271	± 940	± 452	± 564	± 289	± 689	± 447	± 489
$N^{\gamma \rightarrow e}$	4449	4533	14580	6984	9007	5453	13628	9015	10630
$\sigma_{\text{stat.}}$	± 81	± 82	± 148	± 101	± 114	± 89	± 141	± 116	± 124
$N^{h \rightarrow e}$	18004	16690	46087	22846	27170	14766	36533	19793	24650
$\sigma_{\text{stat.}}$	± 173	± 156	± 362	± 223	± 219	± 159	± 294	± 201	± 236

Table 5.7: Top: Summary of extracted $N^{Q \rightarrow e}$ yields in bins of E_T and $|\eta|$. The errors are purely statistical. Bottom: Final $N^{Q \rightarrow e}$ yields in bins of E_T including statistical and systematic uncertainties.

Figure 5.15: Extraction results as a function of E_T (top four panels) and $|\eta|$ (bottom four panels) for the three $Q \rightarrow e$, $\gamma \rightarrow e$, and $h \rightarrow e$ components individually as well as their combination. The simulation is normalized to the data. The data in the top four panels is rescaled to 1.3 pb^{-1} from lower luminosities where necessary to remove discontinuities in the E_T spectrum. Only statistical errors are shown.

Figure 5.16: E_T -distribution of the extracted $Q \rightarrow e$ signal electron component before (left) and after (right) reweighting to 1.3 pb^{-1} . Statistical and systematic uncertainties are shown.

four panels) for the three components individually as well as their combination. The simulation is normalized to the data. The data in the top four panels has been rescaled to 1.3 pb^{-1} from lower luminosities where necessary to remove discontinuities in the E_T spectrum. Only statistical errors are shown in Fig. 5.15. Fig. 5.16 shows the extracted raw (left) and luminosity-corrected (right) $N^{Q \rightarrow e}$ yield including statistical and systematic uncertainties.

5.7 Efficiency Determination and Unfolding Procedure

This section describes the determination of the signal efficiencies associated to the selection requirements of the analysis as well as the migration correction procedure used to unfold the measured spectrum to true electron p_T . While the trigger efficiencies are extracted from data as detailed in section 5.7.1, the efficiencies associated to the reconstruction and identification of electrons as well as the bin migration correction are taken from simulation as further described in section 5.7.4. Input to the final efficiency systematic uncertainties are a number of cross-checks of the efficiencies of the preselection requirements detailed in section 5.7.2 as well as a complete measurement of the identification efficiency as further described in section 5.7.3.

5.7.1 Trigger Efficiency Measurement

The efficiency of the triggers described in section 5.4 to select signal electrons is measured with data generated from independent or lower threshold triggers. For the L1_EM2 and L1_EM5 items the data samples are obtained from a very inclusive minimum bias trigger based on hit information in the minimum bias trigger scintillators. The data sample to measure the L1_EM10 and L1_EM14 efficiencies is generated with the L1_EM5 trigger which is fully efficient in the E_T region where these triggers are used. For comparison the trigger efficiencies are also extracted from simulated signal samples. Fig. 5.17 shows the trigger efficiencies as a function of electron E_T , where good agreement is observed between data and simulated samples. The small residual difference is due to the different fraction of the $h \rightarrow e$ component in the data and simulated samples and is taken as the systematic uncertainty of the measurement. The final values with associated statistical and systematic uncertainties are summarized in Tab. 5.8.

Figure 5.17: Measured efficiencies for the L1_EM2, L1_EM5, L1_EM10 and L1_EM14 trigger items as a function of the E_T -bins of the measurement.

E_T bin [GeV]	7-8	8-10	10-12	12-14
EM2	$99.6^{+0.2}_{-0.3} \pm 0.2$	$100.0^{+0.0}_{-0.2} \pm 0.5$	$100.0^{+0.0}_{-0.6} \pm 0.2$	$98.9^{+0.8}_{-1.6} \pm 1.0$
EM5	-	-	$98.9^{+0.6}_{-1.0} \pm 0.0$	$98.9^{+0.8}_{-1.6} \pm 0.8$
Total	$99.6^{+0.2}_{-0.3} \pm 0.2$	$100.0^{+0.0}_{-0.2} \pm 0.5$	$99.0^{+0.5}_{-1.0} \pm 0.0$	$98.9^{+0.8}_{-1.6} \pm 0.9$

E_T bin [GeV]	14-16	16-18	18-20	20-22	22-26
EM2	$100.0^{+0.0}_{-3.0} \pm 0.0$	$100.0^{+0.0}_{-3.5} \pm 0.0$	$100.0^{+0.0}_{-2.8} \pm 0.0$	$100.0^{+0.0}_{-2.8} \pm 0.0$	$100.0^{+0.0}_{-2.8} \pm 0.0$
EM5	$100.0^{+0.0}_{-3.0} \pm 0.0$	$100.0^{+0.0}_{-3.5} \pm 0.0$	$100.0^{+0.0}_{-4.9} \pm 0.0$	$100.0^{+0.0}_{-5.9} \pm 0.0$	$100.0^{+0.0}_{-5.9} \pm 0.0$
EM10	$93.8^{+0.5}_{-0.5} \pm 0.3$	$97.4^{+0.4}_{-0.5} \pm 1.1$	$99.2^{+0.3}_{-0.3} \pm 0.0$	$99.2^{+0.4}_{-0.5} \pm 0.4$	$99.6^{+0.2}_{-0.3} \pm 0.4$
EM14	-	-	$89.8^{+1.0}_{-1.1} \pm 0.4$	$96.0^{+0.8}_{-1.0} \pm 0.4$	$98.2^{+0.5}_{-0.6} \pm 0.3$
Total	$96.4^{+0.3}_{-1.5} \pm 0.2$	$98.5^{+0.2}_{-1.7} \pm 0.6$	$92.1^{+0.8}_{-1.3} \pm 0.3$	$96.8^{+0.7}_{-1.4} \pm 0.4$	$98.6^{+0.4}_{-1.0} \pm 0.3$

Table 5.8: Trigger efficiency (in percent) for each E_T -bin of the measurement shown per trigger as well as the total luminosity weighted average. Statistical and systematic uncertainties are shown.

5.7.2 Preselection Efficiency

An unbiased measurement of the efficiencies of the preselection requirements listed in section 5.5 from data is not possible due to the low signal purity at this stage of the selection ($\sim 2\%$). The preselection efficiencies used in the cross-section measurement are therefore taken from simulation as further described in section 5.7.4. In this section a number of cross-checks are presented, the results of which are not directly used for the final measurement, but provide some additional confidence in the values obtained from simulation.

The efficiency of the primary vertex requirement is measured from data using events with reconstructed electrons that pass the trigger and data quality requirements. It amounts to 99.84 % averaged across all p_T bins with negligible statistical uncertainty. In simulation the vertex finding probability is 100 %.

The efficiency of the run-dependent fiducial cuts related to the removal of problematic calorimeter regions in η and ϕ is expected to be well modelled in simulation. This is verified in data-taking periods A-C where the modeling of problematic calorimeter regions reflected the actual detector status. An agreement within 0.1 % is observed.

The efficiency of the silicon hits requirement associated to reconstructed electron tracks (≥ 4 hits) can be estimated from the combined $h \rightarrow e$ and $Q \rightarrow e$ components which exhibit similar track characteristics. However, late photon conversions which are present at this stage of the selection are found to bias the results. Therefore the requirement of one additional silicon hit is made to reduce photon conversions to a negligible level. The efficiency of reconstructing four silicon hits is then approximated by the efficiency of measuring five hits given that one has already been found. Agreement within 0.1 % between data and simulation is observed using this technique.

The efficiency of the TRT hits requirement (≥ 10 hits) is measured in a similar fashion requiring at least four silicon hits on the same reconstructed electron track. The observed differences between the measured values and those obtained from simulation range between 0.2-2 %.

E_T bin [GeV]	7-8	8-10	10-12	12-14	14-16	16-18	18-20	20-22	22-26
Silicon hit efficiency [%]									
MC signal	100.0(-)	100.0(-)	100.0(-)	100.0(-)	100.0(-)	99.99(-)	99.99(-)	99.99(-)	100.0(-)
MC	99.88(1)	99.87(1)	99.87(1)	99.85(1)	99.82(2)	99.85(2)	99.87(2)	99.87(3)	99.87(2)
Data	99.86(1)	99.84(1)	99.82(-)	99.82(1)	99.76(1)	99.77(1)	99.79(-)	99.80(1)	99.80(1)
TRT hit efficiency [%]									
MC signal	96.65(2)	97.25(1)	97.70(2)	98.01(3)	98.27(3)	98.41(4)	98.59(5)	98.79(6)	98.86(5)
MC	98.61(2)	98.72(2)	98.84(2)	98.90(3)	98.92(4)	98.94(5)	98.99(6)	99.19(7)	99.09(6)
Data	98.69(2)	98.79(2)	98.81(1)	98.89(2)	99.01(1)	99.06(2)	99.13(1)	99.16(1)	99.26(1)
f_1 cut efficiency [%]									
MC signal	99.67(5)	99.70(4)	99.67(7)	99.7(1)	99.7(1)	99.4(3)	100(-)	100(-)	99.1(5)
MC	96.57(3)	96.68(3)	96.98(5)	97.30(7)	97.46(9)	97.6(1)	98.0(2)	97.9(2)	98.1(2)
Data	96.18(5)	96.25(5)	97.66(2)	97.69(3)	98.12(2)	98.19(3)	98.64(2)	98.68(2)	98.74(2)

Table 5.9: *Efficiency measurements of some preselection requirements with data. “MC signal” is the efficiency for electrons from heavy-flavour decays, while “MC” shows the prediction from the minimum bias simulated sample, with all background contributions included. Statistical errors are shown in brackets where they could be evaluated. Taken from reference [134].*

The efficiency of the f_1 requirement is estimated with a conversion enhanced sample which is constructed by rejecting events with a hit in the pixel B-layer.

The efficiencies of the silicon and TRT hit, as well as the f_1 requirement are summarized in Tab. 5.9 as a function of the E_T bins of the cross-section measurements.

5.7.3 Identification Efficiency

As for preselection, the efficiency of the electron identification criteria described in section 5.5 is taken from simulation. The associated uncertainty, however, is obtained from comparison to a direct data measurement using the so-called Tag & Probe method (T&P) in events where both heavy quarks decay semi-leptonically into electrons. The T&P method is based on the classification of the reconstructed electrons in such di-electron events into a tag electron which fulfills tight selection requirements and one or more probe electrons. The identification efficiency is then evaluated as the number of probes fulfilling the preselection and identification requirements divided by the number of probes fulfilling the preselection requirements only according to

$$\epsilon_{\text{ID}}^{\text{T\&P}} = \frac{N_{\text{probe \& identified}}}{N_{\text{probe}}}. \quad (5.9)$$

Due to the non-resonant character of heavy-flavour production, however, the probe sample does not reach a sufficient purity of signal electrons to directly perform the efficiency calculation shown in equation 5.9. Instead, the $Q \rightarrow e$ component must be extracted from the probe sample in a similar way as it has been done for the signal extraction described in section 5.6.

In the following the T&P selection, the signal extraction method, the evaluation of the associated systematic uncertainties, and the final results are summarized. Further details can be found in reference [134].

Tag & Probe Selection

Events from the full analysis data sample are selected if they contain a tag electron, a probe electron, and fulfill a number of combined selection criteria. The tag electron is required to match the electromagnetic cluster of the corresponding L1 trigger item within $\Delta R < 0.15$ to ensure that it is responsible for triggering the event. In this way a possible trigger bias in the sample of probe electrons is avoided. Tag electrons are further required to have a transverse energy higher than 7 GeV, to be within the pseudorapidity range of the measurement, and to not fall into problematic regions of the calorimeter. In addition, **medium** identification criteria (see Tab. 4.3) augmented by the n_{BL} , n_{TRT} , and f_{TR} requirements listed in Tab. 4.2 are applied.

The selection of probe electrons is identical to the preselection requirements described in section 5.5 with the exception that a lower minimum cut on the transverse energy of $E_T > 3$ GeV is applied such that the efficiency measurement can be extended to this region.

Finally, tag and probe electron pairs are selected if they are oppositely charged and do not fall into the invariant mass window of the J/ψ ($m_{ee} < 2.8$ or $m_{ee} > 3.2$) or Z -boson ($m_{ee} < 70\text{GeV}$) resonances. The opposite charge requirement allows to further increase the signal purity and the invariant mass cuts reduce the contribution from isolated electrons which are not of interest to

this measurement. The impact of isolated electrons originating from the Drell-Yan process which cannot be suppressed with invariant mass requirements has been estimated from simulation and is expected to be small.

After application of the selection requirements described above 12 372 probe electron remain at preselection level and 2 312 after the additional application of the identification criteria.

Signal Extraction

Based on simulation, the fraction of signal electrons in the probe sample amounts to approximately 9 % at preselection level and 31 % after identification requirements. It is thus necessary to apply a signal extraction at both levels of the probe selection to extract the exact number of signal electrons and derive the efficiency.

The technique chosen for the signal extraction is a slightly different implementation of the Tiles Method described in section 5.6. The discriminating variables chosen in this implementation are f_{TR} in four bins (0-0.1, 0.1-0.2, 0.2-0.3, 0.3-1), n_{BL} in two bins (0 and ≥ 1), and f_1 in five bins (0.1-0.2, 0.2-0.3, 0.3-0.4, 0.4-0.5, 0.5-1). This particular choice has been optimized and validated on simulated samples. It has been developed independently from the baseline signal extraction of section 5.6 which explains the different setup.

Due to the limited statistics of the probe samples given the binning in the discriminating variables described above, the measurement was performed in a coarser transverse energy (3-5 GeV, 5-7 GeV, 7-10 GeV, 10-16 GeV, 16-26 GeV) and pseudorapidity (0-0.625, 0.625-1.37, 1.52-2) binning.

Systematic Uncertainties

Systematic uncertainties on the identification efficiency arise from the signal extraction method described above. The procedure used to evaluate the uncertainties is very similar to that used in the baseline implementation of the signal extraction described in section 5.6. The pdfs of the discriminating variables for the $Q \rightarrow e$ and $\gamma \rightarrow e$ components were shifted to model discrepancies between data and simulation. The bias due to correlations in the $h \rightarrow e$ component was evaluated with pseudo experiments. The electron energy scale uncertainty was determined according to the previously described recommendations of the ATLAS electron /

Source	Uncertainty on T&P efficiency [%]
Correlations in the $h \rightarrow e$ component	3.8
f_{TR} mismodeling in the $Q \rightarrow e$ and $\gamma \rightarrow e$ components	2.2
f_1 mismodeling in the $Q \rightarrow e$ and $\gamma \rightarrow e$ components	2.8
n_{BL} mismodeling in the $Q \rightarrow e$ and $\gamma \rightarrow e$ components	4.4
Electron energy scale	2.8
Total uncertainty	7.4

Table 5.10: Summary of relative systematic uncertainties associated to the measurement of the identification efficiency [134].

Figure 5.18: *Left: T&P measured electron identification efficiency with statistical and statistical plus systematic uncertainties as a function of reconstructed E_T . Right: Ratio of measured and true identification efficiency with statistical and statistical plus systematic uncertainties as a function of reconstructed E_T . Taken from [134].*

photon performance group. Some of the uncertainties mentioned above are correlated and may partially cancel in the ratio of the efficiency calculation. To reduce the impact of the limited statistics in the probe samples the uncertainties were determined collectively in the range $3 < E_T < 26$ GeV. The final systematic uncertainties associated to the identification efficiency measurement are listed in Tab. 5.10. The total uncertainty amounts to 7.4 %.

Results

Fig. 5.18 shows the results of the efficiency measurement compared to simulation as a function of E_T (left) as well as the ratio of the measurement and simulation (right). Agreement is observed within the uncertainties. Since a determination of the identification efficiency within the transverse energy bins of the cross-section measurement was not possible the central values for the final result are taken from simulation. The uncertainties, however, are derived from the measurement to account for possible discrepancies in the efficiency between data and simulation. The uncertainties are determined by fitting a constant through the measured E_T points in Fig. 5.18 (right) taking into account their statistical uncertainty only. From this procedure a relative uncertainty of 5.4 % is obtained which is combined with the 7.4 % systematic uncertainty of the measurement to obtain a total uncertainty of 9.2 % on the identification efficiency.

5.7.4 Combined Efficiency and Migration Correction

Based on the cross-checks performed in section 5.7.2 and the measurement in section 5.7.3, the central values of the combined reconstruction and identification efficiencies $\epsilon_{\text{reco+ID}}$ within the E_T range of the cross-section measurement are taken from the high statistics PYTHIA simulated heavy-flavour samples listed in Tab. 5.1. The evolution of the efficiencies after consecutive application of the analysis cuts as a function of true electron p_T is shown in Fig. 5.19 (left). These efficiencies include a matching in ΔR between reconstructed and true generator-level electrons from B - and D -hadron decays within the pseudorapidity region of the measurement. The efficiencies are deteriorated by the non-isolated nature of heavy flavour semi-leptonic decays.

Figure 5.19: *Left: Cumulative efficiency distributions of all selection criteria as a function of true electron p_T obtained from simulation. Right: Combined reconstruction and identification efficiencies before and after bin-to-bin migration corrections including statistical and systematic uncertainties as a function electron p_T . Taken from [134].*

The reconstruction and identification efficiencies are combined with bin migration corrections to unfold the number of electrons in a given reconstructed E_T bin to the number of electrons in the corresponding true p_T bin. Based on the good transverse momentum resolution for electrons from heavy-flavour decays shown in Fig. 5.20 and assuming that the slope of the transverse energy distribution is well-modelled in simulation, a bin-by-bin correction is used to unfold the measured electron spectrum. The corresponding migration correction factor $C_{\text{migration}_i}$ for p_T -bin i is defined as

$$C_{\text{migration}_i} = \frac{N^{Q \rightarrow e}(p_{T_i}^{\text{true}}, E_T^{\text{reco}} > 7 \text{ GeV})}{N^{Q \rightarrow e}(E_{T_i}^{\text{reco}})}. \quad (5.10)$$

The calculation of $C_{\text{migration}_i}$ is based on the high statistics PYTHIA simulated samples (see Tab. 5.1) after application of all selection requirements of the analysis and the matching of true and reconstructed electrons. The performance of the bin-by-bin unfolding procedure is compared to the alternative Singular Value Decomposition (SVD) [160] and Bayesian [161] unfolding methods, the details of which can be found in the precited references. In contrast to equation 5.10, these unfolding techniques necessitate the application of an additional Monte-Carlo based correction factor to take the migrations of true electrons from outside the measurement range ($p_T^{\text{true}} < 7 \text{ GeV}$) into account. This E_T -dependent correction factor is defined as

$$F_i = \frac{N^{Q \rightarrow e}(E_{T_i}^{\text{reco}}, p_{T_i}^{\text{true}} > 7 \text{ GeV})}{N^{Q \rightarrow e}(E_{T_i}^{\text{reco}})}. \quad (5.11)$$

The inputs to the unfolding routine are the F_i -corrected data distribution and a Monte-Carlo based square response matrix describing the probability of reconstructed electrons in a given E_T bin i to migrate to a given true electron p_T bin j . Tab. A.3 in appendix A.2 compares the results of the bin-by-bin unfolding, the bayesian approach with $n = 1$ and $n = 4$ iterations, and the SVD technique with regularization parameters $k = 2$ and $k = 4$. A number of closure

Figure 5.20: *Left: Transverse energy resolution of electrons from heavy-flavour decays obtained from PYTHIA simulated samples. The small non-gaussian tails on both sides of the resolution curve result from energy losses due to Bremsstrahlung at the lower end and from nearby hadronic activity contributing to the electron cluster energy at the higher end. Right: Correlation between true electron p_T and reconstructed electron E_T based on PYTHIA simulated samples. Taken from [134].*

tests are shown for different input samples, as well as for the extracted data themselves. Good agreement within $\sim 2\%$ is observed between the methods. These studies thus support the initial choice of the bin-by-bin unfolding procedure which is preferred also for the simplicity of its implementation.

The combined migration corrected reconstruction and identification efficiencies $\epsilon_{\text{reco+ID}}/C_{\text{migration}}$ are shown in Fig. 5.19 (right). The corresponding values are listed in Tab. 5.11. The systematic uncertainty in each bin is estimated from the difference observed when using simulated samples with an increased amount of material in the ID, in addition to the uncertainty on the identifica-

p_T -bin [GeV]	7-8	8-10	10-12	12-14	14-16	16-18	18-20	20-22	22-26
$\epsilon_{\text{reco}} \pm \sigma_{\text{stat}}$	46.9 ± 0.1	66.9 ± 0.1	73.7 ± 0.2	75.3 ± 0.3	75.9 ± 0.4	76.0 ± 0.5	76.5 ± 0.7	78.3 ± 0.9	77.5 ± 0.9
$\epsilon_{\text{ID}} \pm \sigma_{\text{stat}}$	89.7 ± 0.1	90.1 ± 0.1	88.0 ± 0.2	87.0 ± 0.2	85.1 ± 0.4	83.1 ± 0.5	81.4 ± 0.7	80.0 ± 1.0	77.6 ± 1.0
$\epsilon_{\text{reco+ID}} \pm \sigma_{\text{stat}}$	42.0 ± 0.1	60.3 ± 0.1	64.8 ± 0.2	65.5 ± 0.3	64.7 ± 0.4	63.2 ± 0.6	62.3 ± 0.8	62.3 ± 1.0	60.1 ± 1.0
$\pm \sigma_{\text{syst}} \text{ (d. m.)}$	± 5.4	± 3.0	± 1.4	± 0.9	± 0.4	± 1.6	± 1.5	± 2.4	± 1.3
$C_{\text{migration}} \pm \sigma_{\text{stat}}$	70.6 ± 0.4	95.2 ± 0.5	104.9 ± 0.8	100.8 ± 1.1	97.0 ± 1.5	97.2 ± 1.5	92.8 ± 2.6	94.8 ± 3.4	86.5 ± 3.6
$\pm \sigma_{\text{syst}} \text{ (d. m.)}$	± 3.3	± 3.9	± 6.5	± 6.5	± 4.5	± 5.6	± 4.5	± 6.5	± 9.0
$\frac{\epsilon_{\text{reco+ID}}}{C_{\text{migration}}} \pm \sigma_{\text{stat}}$	59.5 ± 0.4	63.3 ± 0.3	61.8 ± 0.5	64.9 ± 0.8	66.7 ± 1.1	65.0 ± 1.5	67.1 ± 2.1	65.7 ± 2.6	69.5 ± 3.1
$\pm \sigma_{\text{syst}} \text{ (d. m.)}$	± 5.3	± 5.4	± 4.7	± 4.7	± 3.2	± 5.0	± 4.5	± 6.4	± 5.0
$\pm \sigma_{\text{syst}} \text{ (T\&P)}$	± 5.6	± 6.0	± 5.8	± 6.1	± 6.3	± 6.1	± 6.3	± 6.2	± 6.5
Total uncertainty	± 7.7	± 8.1	± 7.5	± 7.7	± 7.1	± 8.0	± 8.0	± 9.2	± 8.8

Table 5.11: *Summary of reconstruction and identification efficiencies as well as migration corrections (in percent) together with the associated statistical and systematic uncertainties in bins of true electron p_T . Systematic uncertainties arising from the distorted material are denoted $\sigma_{\text{syst}} \text{ (d. m.)}$. The further (absolute) uncertainties from the in-situ T\&P identification efficiency measurement and the resulting total uncertainty on $\epsilon_{\text{reco+ID}}/C_{\text{migration}}$ are given in the last two rows.*

Figure 5.21: *Left: Fraction of B -hadron decays to electrons as a function of true electron transverse momentum with the pseudorapidity range of the measurement. A number of different Monte-Carlo generated samples are shown as well as the FONLL prediction and its uncertainty treating the B - and D -hadron rates as correlated, independent, or anti-correlated. Right: Efficiencies and migration correction as well as their ratio for electrons from B - and D -hadron decays as a function of true electron transverse momentum. Taken from [134].*

tion efficiency from the T&P measurement described in the previous section. The uncertainties that result from this procedure cover well the estimated uncertainties of the preselection requirements studied in section 5.7.2.

A series of additional cross-checks is performed on the efficiency and migration correction factors obtained from simulation to evaluate possible theoretical uncertainties. First, the uncertainty on the heavy-flavour composition in the PYTHIA simulated samples is estimated by comparing the B -hadron fraction as a function of the true electron transverse momentum in the measurement range for a number of Monte-Carlo samples as well as the FONLL prediction, where the uncertainties of the $B \rightarrow e$ and $D \rightarrow e$ rates are treated either as fully correlated, independent, or anti-correlated, the latter being the most conservative approach. The results of this comparison are shown in Fig. 5.21 (left). The uncertainty on the B -hadron fraction is then taken from the band that corresponds to the independent variation of the rates. It amounts to 12 % at most inside the measurement range and contains the variations observed with the three generated samples considered. The impact of this uncertainty on the measurement depends on the difference of the combined efficiency and migration correction for electrons from B - and D -hadrons. Fig. 5.21 (right) shows that the individual efficiencies and migration corrections are larger for electrons originating from B -hadron decays while the final ratio of the two quantities agrees within ~ 5 % (relative) for B - and D -hadron decays. The resulting uncertainty related to the B -hadron fraction in the simulated samples is thus only of the order of 1% and is not further propagated into the measurement.

Similarly, the impact of prompt electrons from J/ψ production on the combined efficiency and migration correction was evaluated and found to be of the order of 10 % at most at high transverse momentum. After taking into account the correct mixture of J/ψ events in the analysis sample this effect amounts to less than 1 % difference in $\epsilon_{\text{reco+ID}}/C_{\text{migration}}$ and is neglected for the final measurement.

The default choice of the combined efficiency and migration correction factors obtained from PYTHIA was compared to the values obtained with the lower statistics POWHEG samples listed in Tab. 5.1. Agreement is observed within the statistical uncertainties. The corresponding figures can be found in Fig. A.1 of the appendix A.2.

In addition, the sensitivity of $\epsilon_{\text{reco+ID}}/C_{\text{migration}}$ to pile-up interactions was tested using dedicated samples with pile-up simulation. The impact was found to be negligible.

5.8 Cross-Section Measurement

In this section all elements of the analysis are assembled to the final cross-section measurement. Section 5.8.1 describes the methodology used to calculate the cross-section. The complete list systematic of uncertainties and their treatment in the final cross-section calculation is detailed in section 5.8.2. The final results of the measurement including comparisons to theoretical predictions and complementary measurements are given in section 5.8.3.

5.8.1 Methodology

The differential cross-section for the production of electrons from heavy-flavour decays within the pseudorapidity and transverse energy range of the measurement is determined with equation 5.1. Inputs to the equation are:

- $N_i^{Q \rightarrow e}$, the number of $Q \rightarrow e$ electrons as determined in section 5.6
- $C_{\text{migration}_i}/\epsilon_{(\text{reco+ID})_i}$, the ratio of the migration correction and the combined reconstruction and identification efficiencies as listed in section 5.7.4
- $\epsilon_{\text{trigger}_i}$, the trigger efficiencies determined from data as described in section 5.7.1
- $\int \mathcal{L} dt$, the integrated luminosity of the measurement given in section 5.3
- $\sigma_{\text{accepted}_i}^{W/Z/\gamma^*}$, the accepted cross-section for $W/Z/\gamma^*$ processes obtained from PYTHIA simulated samples after application of all selection criteria

The index i labels the p_T bins of the measurement. For the subtraction of the small $W/Z/\gamma^*$ contribution the W/Z samples are normalized to their NNLO theoretical predictions and the low-mass Drell-Yan contribution to the LO PYTHIA cross-section. The uncertainty of the NNLO W/Z production cross-section calculation [154] of 5 % is applied commonly to the accepted $W/Z/\gamma^*$ cross-section. Since the low mass Drell-Yan contribution is very small this approximation has no significant impact on the final result. An additional uncertainty arises from the reconstruction and identification efficiency for isolated electrons that is contained in the accepted cross-section. This uncertainty is estimated to be 5 % based on T&P efficiency measurements of electrons from W/Z decays [98]. The total impact of these uncertainties on the cross-section measurement is less than 1 %.

5.8.2 Systematic Uncertainties

A full account of the systematic uncertainties associated to the cross-section measurement is given in Tab. 5.12. Some sources of uncertainties are common between the identification efficiency measurement and the signal extraction. These uncertainties are marked with (*) in the table and their uncertainties are propagated to the final measurement taking into account possible correlations.

Shown in Tab. 5.12 are from top to bottom: The uncertainties from correlations in the $h \rightarrow e$ component in the signal extraction and in the efficiency measurement, the effects of the mismodeling of the discriminating variables in the signal extraction and in the efficiency measurement, the electron energy scale uncertainty used in the signal extraction and in the efficiency measurement, the uncertainty related to the possible p_T -dependence of the agreement between the measured efficiency and the expectation from simulation, the material and Monte-Carlo statistical uncertainties on $\epsilon_{\text{reco+ID}}/C_{\text{migration}}$, the Monte-Carlo statistical uncertainties on the $Q \rightarrow e$ and $\gamma \rightarrow e$ pdfs used in the signal extraction, the uncertainties on the trigger efficiencies, the uncertainty on the accepted $W/Z/\gamma^*$ cross-section, and the uncertainty on the integrated luminosity. The total systematic uncertainty on the measured cross-section, taking correlations between the sources into account where necessary, amounts to 15-18 % within the E_T -range of the measurement. This is to be compared with statistical uncertainties from the signal extraction of 2.7-4.3 %.

Source	Uncertainty [%]
Correlations in the $h \rightarrow e$ component	
Signal extraction	7.3
Efficiency measurement	3.8
Mismodeling of discriminating variables	
f_{TR} (*)	4.5
n_{BL} (*)	5.6
E/p	3.2
f_1	2.8
Energy scale (*)	1.5
p_T -dependence of the efficiency measurement	5.5
Material uncertainty on $\epsilon_{\text{reco+ID}}/C_{\text{migration}}$	4.8–9.7
MC statistical error on $\epsilon_{\text{reco+ID}}/C_{\text{migration}}$	0.4–3.5
MC statistical error on templates for signal extraction	0.8–2.5
Trigger efficiency (stat+syst)	< 2
Accepted $W/Z/\gamma^*$ cross-section (MC stat+syst)	< 1
Integrated luminosity	3.4
Total	15–18

Table 5.12: *Summary of systematic uncertainties associated to the cross-section measurement. Sources that are correlated between the signal extraction and the efficiency measurement are marked with (*).*

5.8.3 Results

The final result of the measurement of the differential cross-section for electrons from heavy-flavour decays as a function of p_T in the range 7-26 GeV within the pseudorapidity region $|\eta| < 2.0$ excluding $1.37 < |\eta| < 1.52$ is shown in Fig. 5.22. Within the theoretical and experimental uncertainties the measured cross-section is in good agreement with the FONLL and the POWHEG+PYTHIA predictions. The measured spectrum is observed to fall slightly quicker as a function of p_T . The difference between the NLO+NLL prediction and the NLO-only prediction of FONLL using the same non-perturbative fragmentation function is small but becomes visible towards higher transverse momenta. PYTHIA provides a good description of the shape

Figure 5.22: *Differential cross-section for electron production from heavy flavour decays as a function of electron p_T within $|\eta| < 2.0$, excluding the $1.37 < |\eta| < 1.52$. The ATLAS measurement (dots) with statistical plus systematic errors (blue error bars) is compared to the prediction of the FONLL (light blue band indicating the theoretical uncertainty) and NLO (red line) calculations as well as the predictions of different MC generators (dotted line: PYTHIA, dashed line: POWHEG+PYTHIA, dash-dotted line POWHEG+HERWIG/JIMMY). The ratios of the measured cross-sections to the FONLL central values are shown in the bottom, together with the ratios of the other theoretical predictions to FONLL.*

Figure 5.23: *Left: Differential cross-section for muon production from heavy flavour decays as a function of muon p_T within $|\eta| < 2.5$. The ATLAS measurement (triangles) with statistical plus systematic uncertainties (black error bars) is compared to the prediction of the FONLL (light blue band indicating the theoretical uncertainty) and NLO (red line) calculations as well as the predictions of different MC generators (dotted line: PYTHIA, dashed line: POWHEG+PYTHIA, dash-dotted line POWHEG+HERWIG/JIMMY). The ratios of the measured cross-sections to the FONLL central values are shown in the bottom, together with the ratios of the other theoretical predictions to FONLL. Right: In addition to the information shown in the left panel an NLO-only prediction in combination with non-perturbative fragmentation functions fitted to LEP data using a fixed order NLO calculation is shown. Taken from reference [162].*

of the spectrum but its predicted cross-section is by approximately a factor of two higher than the measurement. A significant underestimation of the cross-section is observed if POWHEG is interfaced to HERWIG and JIMMY for parton showering and the modeling of the underlying event. The difference with respect to the POWHEG+PYTHIA prediction is in part explained by the different decay models as has been further described in section 5.3.

The total cross-section for electron production from heavy-flavour decays integrated over the measurement range, $7 < p_T < 26$ GeV, within the pseudorapidity region $|\eta| < 2.0$, excluding $1.37 < |\eta| < 1.52$, amounts to

$$\sigma_{\text{HF}}^e = 0.946 \pm 0.020(\text{stat.}) \pm 0.146(\text{syst.}) \pm 0.032(\text{lumi.}) \mu\text{b}. \quad (5.12)$$

The total systematic uncertainty on the cross-section has been calculated taking into account possible correlations between E_T bins.

The result shown in Fig. 5.22 is also compared to a similar ATLAS measurement of the cross-section for muon production from heavy-flavour decays [163] within the same transverse mo-

momentum and pseudorapidity range. The two measurements agree within their uncertainties. The total integrated cross-section of the muon measurement within the fiducial region of the electron measurement amounts to

$$\sigma_{\text{HF}}^{\mu} = 0.818 \pm 0.003(\text{stat.}) \pm 0.036(\text{syst.}) \pm 0.028(\text{lumi.}) \mu\text{b}. \quad (5.13)$$

The systematic uncertainties of the muon measurement are smaller compared to the electron measurement due to the lower background contamination which results in smaller systematic uncertainties related to the signal extraction.

The muon cross-section measurement is further extended to the transverse momentum range $4 < p_{\text{T}} < 100$ GeV. The results are shown in Fig. 5.23 (left). It can be seen that at higher transverse momenta the measurement becomes sensitive to the NLL corrections of the FONLL calculation. Fig. 5.23 (right) shows in addition to the default NLO-only prediction which uses the standard FONLL non-perturbative fragmentation functions, an alternative combination of the NLO-only prediction with non-perturbative fragmentation functions fitted to LEP data using a fixed order NLO calculation [162]. Good agreement of this “NLO with NLO NP fit” prediction with the measurement is observed at larger transverse momenta, whereas an underestimation of the data is visible at low and intermediate transverse momenta, which underlines the importance of NLL precision not only at the matrix element level but also in the fit of the non-perturbative fragmentation function.

5.9 Conclusion

The measurement of the differential cross-section for electron production from heavy-flavour decays using the first 1.3 pb^{-1} of data collected by the ATLAS experiment in 2010 at 7 TeV center-of-mass energy in proton-proton collisions was presented in this chapter. The measurement was performed in the electron transverse energy range $7 \text{ GeV} < E_{\text{T}} < 26 \text{ GeV}$ and within pseudorapidities $|\eta| < 2$, excluding $1.37 < |\eta| < 1.52$. The data were collected with L1 electromagnetic triggers and the corresponding trigger efficiencies were determined with independent or lower threshold triggers. After application of electron reconstruction and identification criteria the dominant backgrounds were found to be hadrons faking electrons and electrons from photon conversions. The number of signal electrons was extracted from the data with a three-dimensional extended binned maximum likelihood technique known as Tiles Method. The signal selection efficiencies were derived from simulated samples and the associated systematic uncertainties were determined from independent data measurements and simulated samples assuming additional material with respect to the nominal detector geometry. The measurement was unfolded to true electron transverse momentum using a bin-by-bin unfolding technique. Within the uncertainties, the measured cross-section was found to be compatible with the FONLL prediction at NLO+NLL accuracy. Agreement was also observed with the NLO-only prediction of FONLL using the default non-perturbative fragmentation function obtained from NLO+NLL fits of LEP data. In addition, the measurement is consistent with the POWHEG+PYTHIA prediction, and with a complementary ATLAS measurement with muons in the final state [163]. The total systematic uncertainty of the measurement ranges between 15-18 %, with the largest contribution coming from the determination of the efficiency and migration correction. The measurement was published in [163] together with the complementary inclusive muon analysis.

In conclusion, the measurement provides an important test of the theoretical predictions for heavy-flavour production at the LHC. As one of the first ATLAS measurement at 7 TeV center-of-mass energy, the analysis also constitutes a first evaluation of the electron reconstruction performance in the low electron transverse momentum regime. This is of particular relevance in the context of the search for compressed supersymmetric scenarios described in [chapter 6](#), which is based on events with low- p_T leptons in the final state.

Chapter 6

Search for Compressed Supersymmetry with 2011 Data

This chapter describes the final physics analysis carried out within this thesis - the search for compressed supersymmetric scenarios in final states with missing transverse energy, jets, and one isolated low- p_T lepton using the full ATLAS dataset of $\mathcal{L} = 4.7 \text{ fb}^{-1}$ recorded in 2011 at 7 TeV center-of-mass energy. The physics motivation as well as an overview of the analysis are given in section 6.1. The data and simulated samples of the analysis are listed in section 6.2. The trigger selection, the definition of reconstructed objects and quantities, and the event selection are described in sections 6.3, 6.4, and 6.5 respectively. Sections 6.6 and 6.7 give details about the Standard Model background estimation techniques and the treatment of systematic uncertainties. Section 6.8 describes how all elements of the background estimation and the systematic uncertainties are combined into a global likelihood fit to determine the expected number of background events in the signal regions. The results of the combined fit are summarized in section 6.9. The interpretation of these results in terms of supersymmetric models is described in section 6.10. Finally, section 6.11 concludes on the results of the analysis.

The author's main contributions to the analysis were the selection of triggers, the optimization of the event selection criteria, the production and validation of signal test samples using the MADGRAPH generator for improved modeling of initial state radiation, the customization of an existing fitting framework for the background estimation in the analysis, and the preparation of the final analysis results. The estimation of the multijet background, the final comparison of kinematic distributions between data and simulation, and the development of the fitting software was performed by collaborators, in part working on the two complementary analyses that share the same background estimation framework.

6.1 Introduction

Many searches for the production of supersymmetric particles at the LHC have been carried out to date. Stringent bounds exist on the masses of squarks and gluinos from inclusive searches for an excess of events in final states with missing transverse energy, jets, and additional leptons depending on the supersymmetric scenarios investigated. Fig. 6.1 shows exclusion limits for two particular SUSY models, which were published prior to the work carried out in this thesis

Figure 6.1: Observed (solid red line) and expected (blue dashed line) 95% confidence level exclusion limits for simplified models with gluino (left) and squark (right) production and subsequent decays of the type $\tilde{g} \rightarrow q\bar{q}'\tilde{\chi}_1^\pm \rightarrow q\bar{q}'W^\pm\tilde{\chi}_1^0$ and $\tilde{q}_L \rightarrow q'\tilde{\chi}_1^\pm \rightarrow q'W^\pm\tilde{\chi}_1^0$, respectively. The chargino mass is fixed halfway between the gluino/squark and LSP masses. The color code (right axis) represents the excluded cross-section limit. The results are produced with 1.04 fb^{-1} of ATLAS data taken in 2011 at 7 TeV collision energy using events with missing transverse energy, jets, and leptons with high transverse momentum in the final state. Taken from reference [164].

using the first $\sim 1 \text{ fb}^{-1}$ of ATLAS data collected in 2011. The x -axis in these two figures indicates the mass of the strongly produced supersymmetric sparticles (gluino or squark) while the y -axis denotes the mass of the LSP (neutralino) which stands at the end of the decay chain. Further details on the exact decay topology will follow in the course of this chapter. The red and blue lines in Fig. 6.1 corresponds to the observed and expected limits and the colour code represents the excluded cross-section for a given model point. It can be seen that the analysis is performing well in regions where the mass difference ΔM between the initial sparticles and the LSP is large, while in the so-called compressed mass region with small ΔM close to the diagonal line the sensitivity drops sharply. Given the absence of any experimental evidence for supersymmetric particles in the well-established searches, the previously unexplored compressed mass region could be one of the places where supersymmetric signatures are hidden. The goal of the analysis presented in this thesis is thus to significantly increase the sensitivity for supersymmetric models in this region.

One of the main challenges of this search is the low signal acceptance of compressed SUSY models. Since the predicted supersymmetric particles are very close in mass, only a small amount of kinematic phase-space is available for the decay products which thus have low transverse momentum spectra, typically below the thresholds of unprescaled triggers at the LHC experiments. Due to this soft spectrum, the transverse momentum vectors of the two LSPs in the event are pointing in opposite directions. This results in a cancellation and therefore small missing transverse energy in the event as illustrated in Fig. 6.2 (left). These properties make the events very difficult to discriminate from SM processes.

To obtain an experimental handle on such topologies the analysis performed in this thesis relies on the presence of an initial state radiation (ISR) jet in the event. The system of pair-produced

Figure 6.2: *Illustration of the event topology of a compressed SUSY scenario without (left) and with (right) an initial state radiation (ISR) jet. In the absence of ISR, the pair-produced SUSY particles as well as their decay products are oriented in opposite directions. The transverse momentum vectors of the two LSPs in the event thus cancel approximately leading to a low missing transverse energy signature of the event. If an ISR jet is present, the system of SUSY particles recoils against it such that the two LSPs are more collinear than in the previous scenario and the missing transverse energy is enhanced. Figures are not to scale.*

supersymmetric particles then recoils against this ISR jet as illustrated in Fig. 6.2 (right) such that the transverse momentum vectors of the two LSPs in the event are more aligned and the missing transverse energy signature of the event is enhanced. The acceptance for compressed SUSY scenarios can thus be significantly increased by selecting events with a high transverse momentum ISR jet and large missing transverse energy, while keeping the transverse momentum thresholds of all other objects in the event as low as experimentally possible.

Given the above considerations, the search channel chosen in this analysis is based on events with missing transverse energy, jets, and one isolated low- p_T lepton (hereafter called “soft” lepton) in the final state. At the expense of an acceptance loss for some models, the lepton requirement allows to reject the bulk of multijet backgrounds that would otherwise dominate the analysis.

The soft-lepton analysis is based on so-called simplified models (e.g. [165]). Simplified models consist of simple decay chains with a limited spectrum of hypothetical particles and are parameterized in terms of their production cross-sections, masses, and decay branching ratios. While complete theoretical SUSY models such as CMSSM/MSUGRA often predict very specific mass spectra and event topologies, a diverse selection of simplified models can cover a wider range of signatures and kinematic possibilities. In particular, simplified models allow to scan the full range of supersymmetric particle masses including scenarios with compressed mass spectra.

Figure 6.3: Overview of the common analysis flow of the one-lepton (“soft-lepton” and “hard-lepton”) and “multi-lepton” SUSY searches as published in [166]. The components denoted by grey boxes are not explicitly relevant for the analysis presented in this thesis and are only described as needed for comprehension.

Furthermore, the exclusion limits for simplified models can be reinterpreted in terms of theoretically motivated SUSY models if exclusion limits and acceptance $\times$ selection efficiency maps of the main signal processes exist. The simplified models used in this analysis will be discussed in detail in section 6.2.1.

The soft-lepton analysis is performed together with two complementary searches for SUSY scenarios with larger mass differences between the predicted particles. These analyses are based on events with missing transverse energy, jets, and either one high- p_T lepton (hereafter referred to as “hard-lepton” analysis) or two high- p_T leptons (hereafter referred to as “multi-lepton” analysis). While the two complementary analyses are not the explicit subject of this thesis they are closely connected through the usage of a common analysis framework and common control regions. They will thus be described in this chapter as needed for comprehension.

A schematic overview of the entire analysis flow is shown in Fig. 6.3. As can be seen the analysis contains five exclusive signal regions with different lepton and jet multiplicity requirements that will be detailed in section 6.5. Each signal region is further divided into electron and muon channels, or in the case of the multi-lepton signal regions into electron-electron, electron-muon, and muon-muon channels. The events of the soft-lepton control and signal region are collected with missing transverse energy triggers as detailed in section 6.3. All other control

and signal regions use single-lepton or single-lepton + jet triggers. The determination of the Standard Model backgrounds is described in sections 6.6-6.8. The dominant background sources in the soft-lepton and hard-lepton channels are the production of semi- and fully-leptonic $t\bar{t}$ as well as W +jets events, where the W decays leptonically. For the multi-lepton channels the primary backgrounds are $t\bar{t}$ and Z +jets processes. Minor backgrounds in all signal regions are events from multijet processes, single-top, di-boson, and $t\bar{t}$ plus vector boson production. The background estimation is based on a combined fit as described in section 6.8 using as input dedicated $t\bar{t}$, W +jets, and Z +jets control regions as defined in section 6.6.3, data-driven multijet background estimates as determined in section 6.6.4, and simulated samples for the remaining minor backgrounds as given in section 6.6.5. For an improved agreement with data, the simulated W +jets and Z +jets samples used in the combined background fit are reweighted as a function of vector boson p_T using a separate fit to the Z -boson p_T spectrum as described in section 6.6.2. The associated systematic uncertainties of the analysis are discussed in section 6.7. The results and interpretation in terms of simplified models is given in sections 6.9 and 6.10.

6.2 Data and Simulated Samples

6.2.1 Signal Samples

The simplified models for which the soft-lepton analysis has been conceived are models with gluino or squark pair production followed by a one-step decay into the neutralino via an intermediate chargino according to $\tilde{g} \rightarrow q\bar{q}'\tilde{\chi}_1^\pm \rightarrow q\bar{q}'W^\pm\tilde{\chi}_1^0$ and $\tilde{q}_L \rightarrow q'\tilde{\chi}_1^\pm \rightarrow q'W^\pm\tilde{\chi}_1^0$. The W -boson can be real or virtual and decays to hadronic jets or leptons according to Standard Model branching ratios. The chargino mass is fixed to be mid-way between the mass of the strongly produced gluinos or squarks and the neutralino. This property is generally defined by the mass-ratio parameter

$$x = \frac{m_{\tilde{\chi}^\pm} - m_{\tilde{\chi}_1^0}}{m_{\tilde{g}/\tilde{q}} - m_{\tilde{\chi}_1^0}}, \quad (6.1)$$

where $x = 1/2$ in this particular case. In the models with gluino production, squarks are decoupled by setting their masses to multi-TeV values. Consequently gluinos decay to charginos via three-body decays involving virtual squarks. In the models with squark production only left-handed squarks participate in decays to charginos and only squarks of the first two generations are considered. The free parameters of these simplified models are the masses of the pair-produced squarks or gluinos and the neutralino. In this way regions with compressed mass spectra as well as large- ΔM scenarios are covered.

In addition, though not explicitly constructed for this analysis, a variation of these one-step simplified models is used where the neutralino mass is fixed to 60 GeV and the gluino or squark masses as well as the mass ratio parameter x are varied. In this way scenarios can be covered where the chargino is closer in mass to either the strongly produced initial sparticles or the neutralino. Typical signatures of these models are, depending on x , soft leptons and several high- p_T jets or soft jets and high- p_T leptons in the final state.

The Feynman diagrams of the above described one-step simplified models with gluino and squark production can be found in Fig. 6.4 (a)-(b).

In addition to one-step simplified models, two-step decay scenarios are considered. The multi-lepton analyses have the best expected sensitivity for these models since they are specifically tailored to rich topologies with multiple objects in the final state. The soft-lepton analysis, however, is also expected to add some sensitivity in the compressed mass regions of these models. The production processes for the two-step simplified models are the same as in the one-step case. The subsequent decays can be categorized into three classes.

In the first class of two-step models squarks and gluinos decay via a chargino according to $\tilde{g} \rightarrow q\bar{q}'\tilde{\chi}_1^\pm$ and $\tilde{q}_L \rightarrow q'\tilde{\chi}_1^\pm$. The charginos decay via $\tilde{\chi}_1^\pm \rightarrow \ell\tilde{\nu}_L$ or $\tilde{\chi}_1^\pm \rightarrow \nu\tilde{\ell}_L$ or in the case of third generation sleptons via $\tilde{\chi}_1^\pm \rightarrow \nu\tilde{\tau}_1$. The three generations of sleptons and sneutrinos are assumed to be equally probable.

In the second class of two-step models, decays of the gluinos or left-handed squarks to charginos or neutralinos are implemented according to $\tilde{g} \rightarrow q\bar{q}'\tilde{\chi}_1^\pm$, $\tilde{q}_L \rightarrow q'\tilde{\chi}_1^\pm$, $\tilde{g} \rightarrow q\bar{q}\tilde{\chi}_2^0$, and $\tilde{q}_L \rightarrow q\tilde{\chi}_2^0$. The decays of the pair-produced gluinos or left-handed squarks are configured such that always one chargino and one neutralino is present in the event. The chargino decays proceed as described above and the neutralino decays according to $\tilde{\chi}_2^0 \rightarrow \ell\tilde{\ell}_L$ or $\tilde{\chi}_2^0 \rightarrow \nu\tilde{\nu}$. Again all three generations of sleptons and sneutrinos are allowed with equal probability.

In the third class of two-step models the gluinos and left-handed squarks decay again via $\tilde{g} \rightarrow q\bar{q}'\tilde{\chi}_1^\pm$ or $\tilde{q}_L \rightarrow q'\tilde{\chi}_1^\pm$. The decay of the chargino then proceeds without intermediate sleptons via $\tilde{\chi}_1^\pm \rightarrow W^{(*)\pm}\tilde{\chi}_2^0 \rightarrow W^{(*)\pm}Z^{(*)}\tilde{\chi}_1^0$.

For the first and second class of models the chargino and neutralino masses are fixed mid-way between the gluino/squark and LSP masses. The slepton and neutrino masses, in turn, are set mid-way between the chargino/neutralino and LSP masses. For the third class of models the chargino mass is set midway between the gluino/squark and LSP masses while the mass of the heavier neutralino, $\tilde{\chi}_2^0$, is fixed midway between chargino and LSP.

Representative Feynman diagrams corresponding to the first, second, and third class of two-step simplified models described above are shown in Fig. 6.4 (c)-(d), (e)-(f), and (g)-(h) respectively.

As explained in the previous section of this chapter a key feature of the search for compressed SUSY scenarios is the selection of events with initial state radiation. It is thus important that initial state radiation processes are modeled correctly in the simulated signal samples. Prior to the implementation of the analysis presented here, the simplified model signal samples used by the ATLAS collaboration were generated with HERWIG++ [172]. In HERWIG++ initial state QCD radiation processes are simulated with parton showers (see section 4.4.1). While the parton shower approach provides good modeling in the case of soft and collinear radiation, it is insufficient to describe hard and widely separated partons. In this regime the initial state radiation needs to be calculated at matrix-element level. For a proper description of both regions of phase space, the matrix element and parton shower approaches must be combined with an appropriate matching and removal of overlap.

For this reason the generation of signal samples for the above described simplified models has been moved from the previously used HERWIG++ generator to MADGRAPH5 [173] for the matrix element calculation, interfaced to PYTHIA 6.4.26 [143] for parton showering. The PDF set CTEQ6L1 [150] is used. MLM matching [174] is performed with a scale parameter that is set to one-fourth of the mass of the lightest sparticle in the hard-scattering matrix element following recommendations by the authors of the generator. In the MADGRAPH matrix element calculation of the hard scattering process one additional jet is added to account for initial

Figure 6.5: *Production cross-section as a function of the sparticle mass for $\tilde{g}\text{-}\tilde{g}$ and $\tilde{q}\text{-}\tilde{q}$ production calculated within the MSSM at next-to-leading order including the resummation of soft gluon emission at next-to-leading-logarithmic accuracy (NLO+NLL). Prospino 2.1 and the NLL-fast interpolation tool are used [167–171]. The determination of the cross-section uncertainties is described in section 6.7.9.*

state radiation. The total cross-sections for gluino and squark pair production are calculated in the MSSM at NLO, including the resummation of soft gluon emission at NLL accuracy using Prospino 2.1 and the NLL-fast 1.2 interpolation tool [167–171]. The cross-section and uncertainty values are determined by following where possible the PDF4LHC recommendations [175] as further described in reference [176]. The 68 % confidence level ranges of the CTEQ6.6 [140], CTEQ6.6AS [141] (for α_S -variations), and MSTW2008 [177] PDF sets as well as variations of the renormalization and factorization scales by a factor of two up and down with respect to the PDF sets are used to define an envelope of upper and lower values according to

$$\begin{aligned} U &= \max(\text{CTEQ}_{\text{nom}} + \text{CTEQ}_{\text{up}}, \text{MSTW}_{\text{nom}} + \text{MSTW}_{\text{down}}) \\ L &= \max(\text{CTEQ}_{\text{nom}} - \text{CTEQ}_{\text{up}}, \text{MSTW}_{\text{nom}} - \text{MSTW}_{\text{down}}), \end{aligned} \quad (6.2)$$

where

$$\begin{aligned} \text{CTEQ}_{\text{up}} &= \text{PDF}_{\text{CTEQ,up}} \oplus \text{SCALES}_{\text{CTEQ,up}} \oplus \alpha_{s,\text{up}} \\ \text{CTEQ}_{\text{down}} &= \text{PDF}_{\text{CTEQ,down}} \oplus \text{SCALES}_{\text{CTEQ,down}} \oplus \alpha_{s,\text{down}} \\ \text{MSTW}_{\text{up}} &= \text{PDF}_{\text{MSTW,up}} \oplus \text{SCALES}_{\text{MSTW,up}} \\ \text{MSTW}_{\text{down}} &= \text{PDF}_{\text{MSTW,down}} \oplus \text{SCALE}_{\text{MSTW,down}}. \end{aligned} \quad (6.3)$$

The final cross-section central values are then taken as $\sigma = (U + L)/2$ with an associated symmetric uncertainty of $\Delta_\sigma = (U - L)/2$. The resulting cross-sections and uncertainties for $\tilde{g}\text{-}\tilde{g}$ and $\tilde{q}\text{-}\tilde{q}$ production as a function of the sparticle mass are shown in Fig. 6.5.

Figure 6.6: *Shape comparison between HERWIG++ and MADGRAPH+PYTHIA generated samples at Monte-Carlo truth level for the leading jet p_T (left) and E_T^{miss} distributions (right) for a compressed one-step simplified model with squark production (top) and a large- ΔM model with gluino production (bottom). The transverse momentum spectra of the leading jets and the missing transverse energy distributions are shifted to higher values for the MADGRAPH+PYTHIA samples due to the more realistic modeling of initial state radiation.*

Part of the migration effort to the new generator setup was the production of test samples and a comparison of the relevant kinematic distributions with the well validated HERWIG++ samples used in previous analyses. Fig. 6.6 shows a shape comparison between HERWIG++ 2.5.2 and the MADGRAPH+PYTHIA generated samples for the leading jet p_T (left) and E_T^{miss} distributions (right) for a compressed one-step simplified model with squark production (top) and a large- ΔM model with gluino production (bottom). It can be seen that all expected and desired features of the matrix element + parton shower approach are present. The transverse momentum spectra of the leading jets are harder for the MADGRAPH+PYTHIA samples

Figure 6.7: Comparison of the electron transverse energy (top left), missing transverse energy (top right), transverse mass (bottom left), and missing transverse energy over effective mass distributions in the low electron transverse energy regime using full and fast detector simulation.

which is in part due to the more realistic modeling of initial state radiation in this regime¹. Consequently also the missing transverse energy distribution is shifted to higher values for reasons explained in section 6.1. In addition, the differences between the two generators tend to be larger in the compressed scenarios, in particular in the tail of the missing transverse energy distribution.

Due to the low acceptance for compressed SUSY models it is important to simulate samples in that region with sufficient statistical power. In addition, a large number of model points is required per signal grid to allow for a good coverage over the experimentally accessible mass range of each considered scenario. These requirements quickly result into large and computationally expensive Monte-Carlo production requests, the necessity of which has to be carefully

¹A smaller part of the difference between the two generator setups can be attributed to differences in the parton shower modeling.

Figure 6.8: *Distribution of grid points for the simplified models used in the analysis. The left panel shows the simulated points for the one-step and two-step scenarios with squark and gluino production and fixed mass ratios in the gluino / squark and neutralino mass plane. The blue points indicate the compressed mass region, where the statistical power of some signal samples has been enhanced compared to the large- ΔM region denoted by the red points (see Table 6.1). In the right panel the distribution of simulated model points for the one-step grids with variable x parameter is shown.*

evaluated given the finite resources of the ATLAS Monte-Carlo production system. To keep the CPU processing times within the given constraints, the signal samples have been produced with fast simulation using ATLFAST-II (see section 4.4). While ATLFAST-II is generally well-validated with respect to full simulation, several tests were carried out to check its performance at low lepton- p_T . Fig. 6.7 shows a comparison between full and fast simulation for electron transverse energy and a selection of discriminating variables (see section 6.5.2 for definitions) in the low electron- E_T regime. It can be seen that, within the available statistics, the distributions from full and fast simulation are compatible.

The distribution of grid points for the simplified models used in this analysis is shown in Fig. 6.8. The left panel shows the simulated points for the one-step and two-step scenarios with squark or gluino production and fixed mass ratios with respect to the intermediate particles. In the right panel the distribution of simulated model points for the one-step grids with variable x parameter is shown. Table 6.1 shows the number of events generated per grid point for each simplified model process. If two numbers are listed, the first number refers to the number of events generated in the compressed mass regions (blue points in Fig. 6.8) and the latter to that in the large- ΔM region (red points in Fig. 6.8). Due to the low signal acceptances for compressed scenarios the statistics have been increased for some simplified models in that region.

All samples are passed through the full event simulation chain described in section 4.4. Multiple low- p_T scattering events are overlaid in the simulation to model in-time and out-of-time pile-up interactions (see section 3.1.2).

Physics process		Diagram	Events simulated
$\tilde{g}-\tilde{g} \rightarrow qq\tilde{q}\tilde{q}WW\tilde{\chi}_1^0\tilde{\chi}_1^0$	1-step, $x = 1/2$	Fig. 6.4(a)	80 000 / 20 000
$\tilde{q}-\tilde{q} \rightarrow qqWW\tilde{\chi}_1^0\tilde{\chi}_1^0$	1-step, $x = 1/2$	Fig. 6.4(b)	30 000
$\tilde{g}-\tilde{g} \rightarrow qq\tilde{q}\tilde{q}WW\tilde{\chi}_1^0\tilde{\chi}_1^0, m_{\tilde{\chi}_1^0} = 60 \text{ GeV}$	1-step, variable x	Fig. 6.4(a)	60 000
$\tilde{q}-\tilde{q} \rightarrow qqWW\tilde{\chi}_1^0\tilde{\chi}_1^0, m_{\tilde{\chi}_1^0} = 60 \text{ GeV}$	1-step, variable x	Fig. 6.4(b)	30 000
$\tilde{g}-\tilde{g} \rightarrow qq\tilde{q}\tilde{q}\tilde{\chi}_1^\pm\tilde{\chi}_1^\pm \rightarrow qq\tilde{q}\tilde{q}WZWZ\tilde{\chi}_1^0\tilde{\chi}_1^0$	2-step, via WZ	Fig. 6.4(c)	20 000
$\tilde{g}-\tilde{g} \rightarrow qq\tilde{q}\tilde{q}\tilde{\chi}_1^\pm\tilde{\chi}_1^\pm \rightarrow qq\tilde{q}\tilde{q}ll\nu\nu\tilde{\chi}_1^0\tilde{\chi}_1^0$	2-step, via sleptons	Fig. 6.4(d)	40 000 / 10 000
$\tilde{q}-\tilde{q} \rightarrow qq\tilde{\chi}_1^\pm\tilde{\chi}_1^\pm \rightarrow qqWZWZ\tilde{\chi}_1^0\tilde{\chi}_1^0$	2-step, via WZ	Fig. 6.4(e)	20 000
$\tilde{q}-\tilde{q} \rightarrow qq\tilde{\chi}_1^\pm\tilde{\chi}_1^\pm \rightarrow qqll\nu\nu\tilde{\chi}_1^0\tilde{\chi}_1^0$	2-step, via sleptons	Fig. 6.4(f)	40 000 / 10 000
$\tilde{q}-\tilde{q} \rightarrow qq\tilde{\chi}_1^\pm\tilde{\chi}_2^0 \rightarrow qqlll\nu\tilde{\chi}_1^0\tilde{\chi}_1^0, qql\nu\nu\nu\tilde{\chi}_1^0\tilde{\chi}_1^0$	2-step, via sleptons	Fig. 6.4(g)-(h)	40 000 / 10 000

Table 6.1: Summary of simplified model signal processes and number of generated events per model point for each process. If two numbers are listed, the first number refers to the number of events generated in the compressed mass regions (blue points in Fig. 6.8) and the second number to the number of events in the large- ΔM region (red points in Fig. 6.8).

6.2.2 Background Samples

Table 6.2 lists the simulated samples for the considered Standard Model background processes. The ALPGEN [178] samples are generated with up to five additional partons in the matrix element, except for W + light-flavored jets samples which are generated with up to six partons. For the $Wb\bar{b}$, $Wc\bar{c}$ and Wc cross-sections the LO values from ALPGEN are multiplied with a K-factor of 1.2 based on the K-factor for light-flavored jets. These values are shown for illustrative purposes only. For the final results of this analysis measured cross-section values are used for the W/Z + heavy-flavor-jets samples [179]. The overlap between the heavy-flavored and light-flavored W/Z + jets samples is removed. The single-top cross-sections are taken from MC@NLO and for the s - and t -channels they are listed for a single lepton flavor.

The theoretical cross-sections are calculated with HATHOR 1.2 [180] and the MSTW2008NNLO PDF set for $t\bar{t}$, FEWZ [154] and the MSTW2008NNLO [181] PDF set for W + jets and Z + jets, and MCFM [182] and the MSTW2008NNLO PDF set for di-boson production. The $t\bar{t}+W$ cross-section is taken from reference [183] and the $t\bar{t}+Z$ cross-section corresponds to the leading-order value multiplied by a K-factor of 1.35 derived from the NLO calculation at $\sqrt{s} = 14 \text{ TeV}$ [184].

ALPGEN and MC@NLO are interfaced to HERWIG 6.520 [146] for the modeling of parton showers and fragmentation and to JIMMY 4.31 [147] for underlying event. AcerMC and MADGRAPH are used in combination with PYTHIA 6.4.26. The PDF sets used in the analysis are CTEQ6L1 for the ALPGEN and MADGRAPH samples, CT10 [187] for MC@NLO, and MRSTMCAL (LO**) [188] for HERWIG. For underlying event the ATLAS AUET2B_LO** [189] tunes are used. The ALPGEN and MADGRAPH samples are produced with the MLM matching scheme.

All samples are passed through the full event simulation chain described in section 4.4. Multiple low- p_T scattering events are overlaid in the simulation to model in-time and out-of-time pile-up interactions (see section 3.1.2).

Physics process	Generator	σ [pb]	Accuracy
$t\bar{t}$	ALPGEN 2.13 [178]	166.8	NLO+NLL [180]
$W(\rightarrow \ell\nu) + \text{jets}$	ALPGEN 2.13 [178]	10460	NNLO [154]
$W(\rightarrow \ell\nu) + b\bar{b} + \text{jets}$	ALPGEN 2.13 [178]	130	LO $\times$ K
$W(\rightarrow \ell\nu) + c\bar{c} + \text{jets}$	ALPGEN 2.13 [178]	360	LO $\times$ K
$W(\rightarrow \ell\nu) + c + \text{jets}$	ALPGEN 2.13 [178]	1100	LO $\times$ K
$Z/\gamma^*(\rightarrow \ell\ell) + \text{jets}$ ($m_{\ell\ell} > 40$ GeV)	ALPGEN 2.13 [178]	1070	NNLO [154]
$Z/\gamma^*(\rightarrow \ell\ell) + \text{jets}$ ($10 \text{ GeV} < m_{\ell\ell} < 40 \text{ GeV}$)	ALPGEN 2.13 [178]	3970	NNLO [154]
$Z/\gamma^*(\rightarrow \ell\ell) + b\bar{b} + \text{jets}$ ($m_{\ell\ell} > 40$ GeV)	ALPGEN 2.13 [178]	10.3	LO
Single-top (t -chan)	AcerMC 3.8 [185]	7.0	NLO
Single-top (s -chan)	MC@NLO 4.01 [186]	0.5	NLO
Single-top (Wt -chan)	MC@NLO 4.01 [186]	15.7	NLO
WW	HERWIG 6.520 [146]	44.9	NLO [182]
WZ/γ^* ($m_{Z/\gamma^*} > 60$ GeV)	HERWIG 6.520 [146]	18.5	NLO [182]
$Z/\gamma^*Z/\gamma^*$ ($m_{Z/\gamma^*} > 60$ GeV)	HERWIG 6.520 [146]	5.96	NLO [182]
$t\bar{t}+W$	MADGRAPH5 [173]	0.169	NLO [183]
$t\bar{t}+Z$	MADGRAPH5 [173]	0.120	LO $\times$ K [184]

Table 6.2: List of simulated background event samples used in the analysis with the corresponding Monte-Carlo generators, the production cross-sections, and their accuracies. The notation LO $\times$ K indicates that the process is calculated at LO accuracy corrected by a K-factor derived from the ratio of NLO to LO cross-sections for a closely related process. The $t\bar{t}$, W + light-jets and Z + light-jets samples are normalized using the inclusive cross-sections. The cross-sections for single-top processes in the s - and t -channels as well as for W + light-jets and Z + light-jets processes are listed for a single lepton flavor. Taken from reference [166].

6.2.3 Data Sample

The data used in this analysis were collected from March through October 2011 in 209 detector runs at 7 TeV centre-of-mass energy in proton-proton collisions. Fig. 6.9 (left) shows the cumulative online luminosity distribution versus day during this period as delivered by the LHC (green) and recorded by ATLAS (yellow). The total delivered luminosity amounts to 5.61 fb $^{-1}$ of which 5.25 fb $^{-1}$ were recorded by the ATLAS experiment. The peak instantaneous luminosity reached during that time is $3.65 \times 10^{33} \text{ cm}^{-2} \text{ s}^{-1}$. After application of data quality requirements, the total integrated luminosity used in the analysis amounts to 4.7 fb $^{-1}$ with an uncertainty of 3.9% as determined in [153].

Table 6.3 gives a summary of the 2011 data-taking campaign which is subdivided into a number of periods labelled alphabetically from A to M according to the detector and LHC conditions. During the course of the year the operation parameters evolved to continuously higher luminosities starting with 75 ns proton bunch spacings and only a few colliding bunches at an instantaneous luminosity of $\sim 10^{32} \text{ cm}^{-2} \text{ s}^{-1}$ up to 50 ns bunch spacings and 1380 colliding bunches at an instantaneous luminosity of $3.65 \times 10^{33} \text{ cm}^{-2} \text{ s}^{-1}$. In addition some tests were performed at 25 ns bunch spacing. The maximum luminosity-block averaged number of interactions per beam crossing ranges from approximately 7 up to 19. The luminosity-weighted distribution of the mean number of interaction per bunch crossing is shown in Fig. 6.9 (right). This constitutes a significant increase with respect to the 2010 data-taking (see Fig. 5.4), and requires tightened selection requirements for the reconstructed objects used in the analysis in order to reduce the impact of pile-up interactions. Another note-worthy feature is the loss

Period	Date Range	Run Range	Stable / Ready Lum. [pb ⁻¹]	Max. Inst. Lum. [10 ³⁰ cm ⁻² s ⁻¹]	$\mu_{\text{max}}^{\text{avg}}$	Fill Range	Bunch Count	Bunch Δt [ns]
A	Mar-13 - Mar-21	177531 - 177965	9 / 8.7 (97%)	158	7.3	1613 - 1640	3 - 138	75 - 2500
B	Mar-21 - Mar-24	177986 - 178109	18 / 18 (97%)	249	9.2	1642 - 1647	138 - 194	75
D	Apr-14 - Apr-29	179710 - 180481	188 / 185 (98%)	719	7.6	1710 - 1746	214 - 598	50
E	Apr-30 - May-03	180614 - 180776	54 / 52 (98%)	841	7.8	1748 - 1756	598 - 700	50
F	May-15 - May-25	182013 - 182519	162 / 158 (97%)	1113	8.1	1783 - 1807	14 - 874	50 - 1250
G	May-27 - Jun-14	182726 - 183462	580 / 572 (99%)	1278	8	1809 - 1868	874 - 1042	50
H	Jun-16 - Jun-28	183544 - 184169	291 / 286 (98%)	1278	6.9	1870 - 1902	1041 - 1318	50
I	Jul-13 - Jul-29	185353 - 186493	423 / 411 (97%)	1913	9.2	1936 - 1991	2 - 1331	50 - 2500
J	Jul-30 - Aug-04	186516 - 186755	244 / 240 (98%)	2023	9.8	1992 - 2002	1317	50
K	Aug-04 - Aug-22	186873 - 187815	695 / 684 (98%)	2356	11	2002 - 2040	1317	50
L	Sep-07 - Oct-05	188902 - 190343	1653 / 1621 (98%)	3295	16	2083 - 2183	250 - 1318	50
M	Oct-06 - Oct-30	190503 - 191933	1202 / 1174 (98%)	3896	19	2185 - 2268	2 - 1380	25 - 8000
Total	Mar-13 - Oct-30	177531 - 191933	5521 / 5408 (98%)	3896	19	1613 - 2268	2 - 1380	25 - 8000

Table 6.3: Summary of run periods with stable beams in proton-proton collisions at 7 TeV center-of-mass energy during the 2011 data-taking campaign. For each period important ATLAS detector and LHC parameters are given. “Stable luminosity” refers to the integrated luminosity over all stable beam luminosity blocks while “ready luminosity” corresponds to the subset of stable luminosity where all ATLAS detector systems were at their operational settings for physics data-taking. The difference to the recorded luminosity value in Fig. 6.9 is in part explained by the trigger deadtimes at L1. The quantity $\mu_{\text{max}}^{\text{avg}}$ is the maximum of the luminosity block averaged values of the number of interactions per beam crossing in a given period. The column “bunch count” refers to the number of colliding bunches in the LHC. “Bunch Δt ” is the minimum time between colliding bunches. Period C is omitted since the center-of-mass energy of the LHC during this period was not at its nominal value. The values and ranges in this table are for illustrative purpose only and do not necessarily correspond to those during physics data-taking.

Figure 6.9: *Left: Cumulative distribution of luminosity versus day delivered by LHC (green) and recorded by ATLAS (yellow) during stable beams for proton-proton collisions at 7 TeV centre-of-mass energy in 2011. Right: Luminosity-weighted distribution of the mean number of interactions per bunch crossing for the 2011 dataset. Further details can be found in [153].*

of six LAr Calorimeter FEBs (see section 3.2.2) on 30th April corresponding to 0.4% of the calorimeter cells. The affected region was located at $0.0 < |\eta| < 1.4$ and $-0.74 < \phi < -0.64$ and resulted in decreased efficiency and precision of the reconstructed objects in that region. Four of the six FEBs could be recovered during a LHC technical stop in July and were again available from period I.

6.3 Trigger Selection

To select events with low- p_T leptons in the final state, the primary triggers in this analysis are based solely on missing transverse energy. Single-lepton triggers are used for the complementary high lepton- p_T search channels, which are not the main subject of this thesis. However, since control regions are shared between the analyses, the single-electron and single-muon triggers are also relevant in the context of the soft-lepton analysis.

An overview of the employed triggers as a function of the data taking period is shown in Table 6.4. To adapt to the increasing instantaneous luminosity provided by the LHC throughout the 2011 data-taking campaign, several triggers have been replaced by new triggers with tighter selection requirements.

The missing transverse energy trigger `EF_xe60_noMu` used in periods B-I is seeded by 40 GeV threshold triggers at L1 and L2 (`L1_XE40` and `L2_xe40_noMu`). The trigger `EF_xe60_tight_noMu` used in periods J-K is tightened at L2 with a threshold of 45 GeV (`L2_xe45_noMu`). In periods L-M the trigger `EF_xe60_verytight_noMu` was deployed, which is tightened to 50 GeV at L1 and 55 GeV at L2 (`L1_XE50` and `L2_xe55_noMu`). Fig. 6.10 shows the trigger efficiencies of the `EF_xe60_tight_noMu` (left) and `EF_xe60_verytight_noMu` (right) triggers as a function of offline reconstructed missing transverse energy for data in the corresponding data-taking periods and a simulated one-step simplified signal model in the compressed mass region. The trigger efficiency is determined from events triggered with the `EF_mu18_L1J10` trigger and requiring at

	Period	Trigger Chain
E_T^{miss} triggers	B-I	EF_xe60_noMu (seeded by L1_XE40 $\rightarrow$ L2_xe40_noMu)
	J-K	EF_xe60_tight_noMu (seeded by L1_XE40 $\rightarrow$ L2_xe45_noMu)
	L-M	EF_xe60_verytight_noMu (seeded by L1_XE50 $\rightarrow$ L2_xe55_noMu)
Electron triggers	B-J	EF_e20_medium (seeded by L1_EM14 $\rightarrow$ L2_e20_medium)
	K	EF_e22_medium (seeded by L1_EM16 $\rightarrow$ L2_e22_medium)
	L-M	EF_e22vh_medium1 (seeded by L1_EM16VH and L2_e22vh_medium1) EF_e45_medium1 (seeded by L1_EM30 $\rightarrow$ L2_e45_medium1)
Muons triggers	B-J	EF_mu18 (seeded by L1_MU10 $\rightarrow$ L2_mu18)
	K-M	EF_mu18_L1J10 (seeded by L1_MU10_J10 $\rightarrow$ L2_mu18_L1J10)

Table 6.4: Summary of triggers chains used in the analysis as a function of the data-taking period. The E_T^{miss} -triggers are the primary triggers of the analysis, the electron and muon triggers are used to select events for common control regions with the hard-lepton search channels.

least one reconstructed signal muon (to be defined in section 6.4.2) in the event. The muon term (see section 6.4.4) is removed from the reconstructed missing transverse energy since no muon information is used in the above triggers. It can be seen that the EF_xe60_verytight_noMu trigger and thus also its looser versions are highly efficient at values greater than 180 GeV in missing transverse energy. Moreover, very good agreement with simulated samples is observed in this region. This allows to safely place a lower missing transverse energy requirement of

Figure 6.10: Trigger efficiencies of the EF_xe60_tight_noMu (left) and EF_xe60_verytight_noMu (right) triggers as a function of offline reconstructed missing transverse energy for data and a simulated one-step simplified signal model in the compressed mass region ($\tilde{g}\text{-}\tilde{g} \rightarrow qq\bar{q}\bar{q}WW\tilde{\chi}_1^0\tilde{\chi}_1^0$, $x = 1/2$, $m_{\tilde{g}} = 505$ GeV $m_{\tilde{\chi}_1^0} = 425$ GeV). The muon term is removed from the missing transverse energy as no muon information is used in the above triggers. Both triggers are highly efficient at values greater than 180 GeV in missing transverse energy, where agreement between data and the simulated sample is very good. The discrepancy between data and simulation in the turn-on region is not relevant for this analysis since no events are selected in this region.

Figure 6.11: Trigger efficiencies as a function of transverse lepton momentum for single electron (left) and muon (right) triggers as indicated in the plots. The muon trigger efficiency is shown in the range $|\eta| > 1.05$ only. Taken from references [93] and [191].

180 GeV in the analysis. Further information on the performance of missing transverse energy triggers can be found in reference [190].

The single-electron triggers used in the analysis are `EF_e20_medium` seeded by `L1_EM14` and `L2_e20_medium` for periods B-I, `EF_e22_medium` seeded by `L1_EM16` and `L2_e22_medium` for period K, and a logical OR between triggers `EF_e22vh_medium1` and `EF_e45_medium1`, seeded by `L1_EM16VH` and `L2_e22vh_medium1`, and `L1_EM30` and `L2_e45_medium1` respectively for periods L-M. The letters VH stand for triggers with reoptimized η -dependent transverse energy thresholds and hadronic leakage requirements at L1. The quality criteria `medium` and `medium1` are described in section 4.1.2. The logical OR between the two triggers in periods L-M is used to avoid efficiency losses at high transverse momentum. The efficiency of the single-electron triggers as a function of reconstructed electron transverse momentum is shown in Fig. 6.11. It can be seen that the efficiency rises slowly and reaches its plateau at $p_T \approx 35 - 40$ GeV. In the combined control regions of the analysis, events triggered by electrons with $p_T > 25$ GeV are selected. For an improved description of the data, trigger efficiency scale factors are applied to simulated samples. These scale factors are obtained from comparisons of the trigger efficiency to simulated samples. More information about the performance of electron triggers can be found in [93].

The muon triggers used in the analysis are `EF_mu18` seeded by `L1_MU10` and `L2_mu18` in periods B-J, and the combined muon-jet trigger `EF_mu18_L1J10` seeded by `L1_MU10_J10` and `L2_mu18_L1J10` in periods K-M. As an example, the trigger efficiency for `EF_mu18` is shown in Fig. 6.11 (right) for the pseudorapidity range $|\eta| > 1.05$. In the combined control regions of the analysis events triggered by muons with $p_T > 20$ GeV are selected. Scale-factors are applied to bring simulated samples and data in agreement. More detailed information on the performance of muon triggers can be found in reference [191].

6.4 Definition of Reconstructed Objects and Quantities

The reconstructed objects (electrons, muons, jets) and quantities (missing transverse energy) used in the analysis are defined in this section. All reconstructed objects are defined at preselection level and signal level. The preselected objects with looser requirements are used to veto additional objects (see section 6.5.2) or for background estimation (see section 6.6.4).

6.4.1 Electrons

Electrons are reconstructed with the standard cluster-based reconstruction algorithm described in section 4.3.1. Electrons are preselected if their transverse momentum is within $7 \text{ GeV} < E_T < 25 \text{ GeV}$ in the soft-lepton channel and $E_T > 10 \text{ GeV}$ in the hard-lepton and multi-lepton channels. The pseudo-rapidity range is $|\eta^{\text{clust}}| < 2.47$. In the soft-lepton channel electrons within the transition region $1.37 < |\eta^{\text{clust}}| < 1.52$ between the barrel and endcap calorimeters are excluded. Furthermore, preselected electrons are required to fulfill the **medium++** identification criteria shown in Tab. 4.3. In the case of overlap between an electron and a jet within $\Delta R(\text{electron}, \text{jet}) < 0.2$ the jet is removed (see section 6.4.3). In addition electrons must be isolated from jets within $\Delta R(\text{electron}, \text{jet}) > 0.4$.

Signal electrons are required to fall into the transverse momentum range $7 \text{ GeV} < E_T < 25 \text{ GeV}$ for the soft-lepton channel and $E_T > 25 \text{ GeV}$ for the hard-lepton and multi-lepton channels. In addition to the preselection criteria, **tight++** identification cuts (see section Tab. 4.3) must be fulfilled. An isolation requirement is imposed by selecting only electrons with $p_T^{\text{cone20}}/E_T < 0.10$, where p_T^{cone20} is the sum of the transverse momenta of all tracks in a cone of $\Delta R = 0.2$ around the electron track.

Scale-factors are applied to simulated samples to correct for differences in the electron reconstruction and identification efficiencies between data and simulation. These scale-factors are determined from efficiency measurements based on the previously introduced tag & probe method (see section 5.7). The electron reconstruction efficiencies are determined using $Z \rightarrow ee$

Preselected Electron	
Algorithm	Standard electron reconstruction (see section 4.3.1)
Acceptance	$7 \text{ GeV} < E_T < 25 \text{ GeV}$, $ \eta^{\text{clust}} < 2.47$, $1.37 < \eta^{\text{clust}} < 1.52$ excluded (soft-lepton channels) $E_T > 10 \text{ GeV}$, $ \eta^{\text{clust}} < 2.47$ (hard-lepton channels)
Quality	medium++
Overlap	$\Delta R(\text{electron}, \text{jet}) > 0.4$
Signal Electron	
Acceptance	$7 \text{ GeV} < E_T < 25 \text{ GeV}$ (soft-lepton channels) $E_T > 25 \text{ GeV}$ (hard-lepton channels)
Quality	tight++
Isolation	$p_T^{\text{cone20}}/E_T < 0.10$

Table 6.5: *Summary of the electron selection criteria. The signal selection requirements are applied on top of the preselection.*

Figure 6.12: *Left: Identification efficiency measurement and comparison to simulation for loose++ electrons with transverse energy within 7-50 GeV. The different measurement with $Z \rightarrow ee$, $W \rightarrow e\nu$, and $J/\psi \rightarrow ee$ events are shown. Taken from reference [89]. Right: Scale-factors for tight++ electrons in bins of the transverse electron energy separated by the dashed vertical lines. The values in bins of pseudorapidity (negative to positive from left to right) are given for each transverse energy bin. The scale-factors are applied to the simulated samples.*

events. The electron identification efficiencies at low transverse momentum (up to 20 GeV) are determined with $J/\psi \rightarrow ee$ events. At higher electron- p_T (> 15 GeV) $W \rightarrow e\nu$ and $Z \rightarrow ee$ are available. The identification efficiencies from these measurements for the example of **loose++** identification criteria are shown in Fig. 6.12. The uncertainties are bigger at low transverse momentum due to the more difficult nature of the measurement in this region and reduced statistics due to the usage of heavily prescaled J/ψ triggers. The scale-factors and associated uncertainties, as applied to **tight++** electrons in the analysis, are shown in Fig. 6.12 (right) as a function of electron transverse energy and pseudorapidity.

6.4.2 Muons

“Combined” or “segment-tagged” muons from the **STACO** collection are used in the analysis. This includes muons reconstructed with the **STACO** and **Mutag** algorithms (see section 4.3.2). The acceptance region for preselected muons is $6 \text{ GeV} < p_T < 20 \text{ GeV}$ and $|\eta| < 2.4$ in the soft-lepton channel and $p_T > 10 \text{ GeV}$, $|\eta| < 2.4$ in the hard-lepton and multi-lepton channels. Loose quality criteria are applied. In the ID, the sum of the numbers of hits in the pixel detectors and the number of crossed dead pixel sensors is required to be greater than one. Similarly, the sum of the number of hits in the SCT and the number crossed dead SCT sensors is required to be greater than six. The number of inactive regions (“holes”) in the SCT and pixel detectors traversed by the muon track is required to be smaller than three. In addition, muon candidates are selected if they have at least one hit in the B-layer when it can be expected. In the MS, tracks in the region $|\eta| < 1.9$ are required to have $n_{\text{TRT}} \geq 5$ hits in the TRT, where n_{TRT} is defined as the sum of the number of regular hits along and track and outlier hits. The number of outlier hits is required to be less than $0.9 n_{\text{TRT}}$. For tracks with $|\eta| \geq 1.9$ the same selection is applied only if $n_{\text{TRT}} \geq 5$ hits are found. Otherwise no requirement on the number of TRT hits is applied, since the corresponding muon track could be outside the coverage region of the TRT. Furthermore, preselected muons must be isolated from jets according to $\Delta R(\mu, \text{jet}) > 0.4$.

	Preselected Muon
Algorithm	STACO, combined or segment-tagged muon
Acceptance	$6 \text{ GeV} < p_T < 20 \text{ GeV}$, $ \eta < 2.4$ (soft-lepton channels) $p_T > 10 \text{ GeV}$, $ \eta < 2.4$ (hard-lepton channels)
Quality	Loose
Inner Detector	Number of pixel hits + number of crossed dead pixel sensors ≥ 2 Number of SCT hits + number of crossed dead SCT sensors ≥ 6 Number of pixel holes + number of SCT holes ≤ 3 Number of B-layer hits when a hit is expected ≥ 1
Track Quality	If $ \eta < 1.9$: $n_{\text{TRT}} > 5$ and $n_{\text{TRT}}^{\text{outliers}} < 0.9 n_{\text{TRT}}$ If $ \eta \geq 1.9$ and $n_{\text{TRT}} > 5$: $n_{\text{TRT}}^{\text{outliers}} < 0.9 n_{\text{TRT}}$
Overlap	$\Delta R(\mu, \text{jet}) > 0.4$
	Signal Muon
Acceptance	$6 \text{ GeV} < p_T < 20 \text{ GeV}$ (soft-lepton channels) $p_T > 20 \text{ GeV}$ (hard-lepton channels)
Isolation	$p_T^{\text{cone20}} < 1.8 \text{ GeV}$

Table 6.6: Cuts defining a good-quality reconstructed muon. In this table, n_{TRT} is the total number of TRT hits, including outliers. The signal selection requirements are applied on top of the preselection.

Figure 6.13: Muon reconstruction efficiencies as a function of p_T (left) and $q \times \eta$ (right) obtained from tag $\mathcal{E}$ probe measurements using $J/\psi \rightarrow \mu\mu$ events. The variable q denotes the sign of the charge of the muon. “CB Chain 1” corresponds to combined muons and “CB+ST Chain 1” to combined and segment-tagged muons reconstructed with the STACO chain. Taken from [88].

In addition to the requirements described above, signal muons must be reconstructed within $6 \text{ GeV} < p_T < 20 \text{ GeV}$ in the soft-lepton channel and within $p_T > 20 \text{ GeV}$ in the complementary hard-lepton channel and multi-lepton channels. Finally the sum of transverse track momenta within a cone of $\Delta R = 0.2$ around the muon track, excluding the muon track, must be less than 1.8 GeV .

Scale-factors provided by the ATLAS muon combined performance group are applied to simulated samples in the analysis to correct for differences in the muon reconstruction efficiencies between data and simulation. As for electrons the efficiencies are measured using the tag & probe method. $Z \rightarrow \mu\mu$ events are used over the entire transverse momentum range. The results are found to be in good agreement with complementary measurements using $J/\psi \rightarrow \mu\mu$ events. Fig. 6.13 shows the reconstruction efficiencies measurement using the $J/\psi \rightarrow \mu\mu$ measurement.

6.4.3 Jets

Jets are reconstructed with the **AntiKt4Topo** algorithm and are required to lie within the acceptance region $p_T > 20 \text{ GeV}$, $|\eta| < 2.5$. The EM + JES calibration scheme is applied. Details on the jet reconstruction and calibration in ATLAS are given in section 4.3.3. To avoid double-counting, jets are removed if they fall within a cone of $\Delta R = 0.2$ of an electron as already described above.

Furthermore jet cleaning cuts are applied to reject events with jets classified according to the ATLAS internal scheme “looser bad jets”. This scheme implements two sets of cuts, one to reject jets related to detector problems, and one to discriminate against cosmic ray and non-collision backgrounds. Such non-collision backgrounds can originate from interactions of the incoming protons with residual gas inside the beam pipe (“beam-gas events”) or with the beam collimators in the forward region (“beam-halo events”). The first class of cuts is comprised of requirements on the maximum energy fraction in the hadronic endcap calorimeter (HEC), the calorimeter pulse-shape quality, and negative cell energies to suppress sporadic noise bursts. Similar selections are also applied to reduce rare coherent noise in the electromagnetic calorimeter. The second class of cuts consists of requirements on the fraction of energy deposited in the electromagnetic calorimeters compared to the total energy, the fraction of energy deposited in a single calorimeter layer, and the jet charged fraction, defined as the scalar sum of transverse track momenta associated with the jet divided by the total jet p_T . This set of cuts allows to reject events with jets that do not originate from regions close to the interaction point as would be the case for non-collision backgrounds.

Due to the high pile-up environment with many interactions per bunch crossings it is necessary to introduce dedicated cuts to suppress events with jets that do not originate from the interaction vertex of interest. This is done with a selection on the so-called jet vertex fraction variable $JVF > 0.75$, which is the fraction of the scalar sum of the transverse momenta of tracks in a given calorimeter jet originating from a common primary vertex. Fig. 6.14 (left) shows the sensitivity of the average number of observed jets as a function of the number of primary vertices in $Z \rightarrow ll$ events with and without JVF selection applied. Without pile-up suppression the number of jets increases as expected with the number of primary vertices, an effect that is overestimated in simulation. With pile-up suppression applied the number of jets is approximately constant as a function of primary vertices and in agreement with simulation. The JVF selection is thus very effective in rejecting pile-up jets.

Preselected Jet	
Algorithm	AntiKt4Topo (see section 4.3.3)
Acceptance	$p_T > 20 \text{ GeV}$, $ \eta < 2.5$
Quality	reject events with looser bad jets
Jet vertex fraction	JVF > 0.75
Overlap	$\Delta R(\text{jet}, \text{electron}) > 0.2$
Signal Jet	
Acceptance	$p_T > 25 \text{ GeV}$, $ \eta < 2.5$
b-jets	JetFitterCombNN > 1.8

Table 6.7: *Summary of the jet and b-jet selection criteria. The signal selection requirements are applied on top of the preselection. The b-tagging definition is given in the text.*

In addition to the preselection criteria described above, signal jets are required to have a transverse momentum of 25 GeV. Signal b-jets are identified with the **JetFitterCombNN** tagging algorithm, which is based on a neural network approach using the output weights of the **JetFitter** and **IP3D** algorithms as described in [192]. The **JetFitter** algorithm [193] implements pattern recognition using a Kalman filter (e.g. [194]) to find the flight path of the b- or c-hadron between the primary and secondary vertices. The b-, c-, and light jets are separated with a likelihood approach using the decay length significance and vertex properties. The **IP3D** algorithm is an impact parameter based algorithm which builds a likelihood ratio based on the comparison of two-dimensional histograms of the transverse and longitudinal impact parameter significances from simulation with the measured quantities.

The operating point chosen for the **JetFitterCombNN** algorithm is at 60 % efficiency and translates to a requirement on the corresponding output variable to be larger than 1.8. Small jet- p_T dependent differences between the measured b-tagging and mistagging² efficiencies and those obtained from simulation are observed and correction factors as determined by the ATLAS heavy flavour tagging group are applied to simulated samples.

6.4.4 Missing Transverse Energy

The missing transverse energy used in the analysis is calculated with the **Simplified20_RefFinal** algorithm. **Simplified20_RefFinal** is an object-based missing transverse energy quantity, calculated as the modulus of the negative vectorial sum of the reconstructed objects found in the event. The considered objects are jets reconstructed with the **AntiKt4Topo** algorithm calibrated at EM+JES with transverse momentum greater than 20 GeV, medium electrons (see table 4.3) with transverse momentum greater than 10 GeV, and muons from the **STACO** collection which are separated from the jets described above by $\Delta R > 0.3$. All topological calorimeter clusters not belonging to any of the objects mentioned above are combined into a so-called “cell-out” term and are added to the missing transverse energy. Since the electrons and muons used in the analysis are not known a priori in the missing transverse energy calculation during event reconstruction the corresponding terms are recalculated using leptons according to the object definitions described in the previous sections. The observed missing transverse energy resolution in 2011 for $Z \rightarrow ee$ and $Z \rightarrow \mu\mu$ events is shown in Fig. 6.14 (right).

²The mistagging efficiency is defined as the number of light jets tagged as b-jets over the total number of light jets.

Figure 6.14: *Left: Average number of jets as a function of the number of primary vertices in $Z \rightarrow ll$ events with and without JVF selection applied. It can be seen that the JVF selection is highly effective in rejecting events with pile-up jets. Right: Resolution of missing transverse energy as a function of the total transverse energy measured in $Z \rightarrow ee$ and $Z \rightarrow \mu\mu$ events. Taken from reference [90].*

6.5 Event Selection

This section describes the preselection cuts applied to data and simulated samples as well as the optimization procedure and final signal region selections with event cut-flow tables and comparisons between data and simulation for the most relevant distributions.

6.5.1 Preselection

Events in data are selected if the detector runs and luminosity blocks in which they are contained are included in the good run list (GRL) for physics analysis. This GRL requires all detector components to be in a state that is deemed acceptable for physics analysis by the ATLAS data quality groups. Furthermore a primary vertex in the nominal beam spot region with at least five associated tracks is required. In case more than one vertex fulfilling these criteria is found the one with the highest summed $|p_T^2|$ of associated tracks is chosen. In addition to the jet cleaning described in section 6.4.3, a global LAr calorimeter noise burst flag is checked. To mitigate the effects of the six lost LAr Calorimeter FEBs (see section 6.2.3), which affect approximately 20 % of the data, a so-called “smart event-veto” is implemented. If a jet falls into a dead region its transverse energy is estimated from energy depositions in adjacent cells. If the estimated jet transverse energy projected along the axis of the missing transverse energy vector is higher than 10 GeV and constitutes more than 10 % of the missing transverse energy the event is rejected. Events with electrons that fall into the problematic calorimeter region are rejected. Further, events with “badly” reconstructed muons $\sigma(q/p)/|(q/p)| > 0.2$, where p denotes the momentum and q the sign of the electric charge, are rejected. Events with cosmic muons are vetoed by requiring $|z_0^{\text{PV}}| < 1 \text{ mm}$ and $|d_0^{\text{PV}}| < 0.2 \text{ mm}$ on the longitudinal and transverse impact parameters of the muon tracks with respect to the primary vertex.

6.5.2 Signal Region Optimization

As demonstrated in the introduction of this chapter (section 6.1), the typical experimental signature of compressed SUSY scenarios, consists of low- p_T objects and low missing transverse energy due to oppositely aligned transverse momentum vectors of the LSPs. This makes it difficult to trigger on such events and to discriminate them against Standard Model backgrounds. It has been demonstrated, however, that the sensitivity to such scenarios can be significantly improved if events with hard initial state radiation jets and thus larger missing transverse energy are selected.

The event selection of this analysis is thus based on the requirement of at least one energetic jet, significant missing transverse energy, as well as low transverse momentum secondary jets, and soft electrons or muons. In addition to the selection criteria associated to the multiplicities and transverse momenta of the reconstructed objects, the following discriminating variables are used:

The transverse mass, is computed from the lepton transverse momentum p_T^ℓ and the missing transverse momentum $\vec{p}_T^{\text{miss}}$ in the event according to:

$$m_T = \sqrt{2p_T^\ell E_T^{\text{miss}}(1 - \cos(\Delta\phi(\vec{\ell}, \vec{p}_T^{\text{miss}})))}, \quad (6.4)$$

For the dominant Standard Model backgrounds, the m_T distribution has an endpoint around the value of the W boson mass thus providing separation power between Standard Model and beyond Standard Model events. Another discriminating variable is the ratio of $E_T^{\text{miss}}/m_{\text{eff}}$, where m_{eff} is defined as the scalar sum of the transverse momenta of the leptons and the n leading jets in the event, where n is the minimum jet multiplicity required in the analysis. The related quantity $m_{\text{eff}}^{\text{inc}}$ is defined as m_{eff} but makes use of all jets in the event. It is only used in the hard-lepton and multi-lepton analyses. In addition, angular variables between the reconstructed objects were considered but did not provide enough discriminating power to justify their usage.

In order to determine an optimal set of selection requirements the cuts on the observables described above were varied using the following values:

- $E_T^{\text{miss}} > [220, 250, 270]$ GeV
- $E_T^{\text{miss}}/m_{\text{eff}} > [0.2, 0.3, 0.4]$
- Leading jet $p_T^{\text{jet},1} > [100, 130, 160]$ GeV
- Jet multiplicity $N_{\text{jet}} \geq [2, 3]$

The selection criteria on the remaining variables were fixed. The lepton transverse momentum range was set to $7 \text{ GeV} < E_T < 25 \text{ GeV}$ for electrons and $6 \text{ GeV} < p_T < 20 \text{ GeV}$ for muons as described in section 6.4. The lower thresholds are governed by the availability of reliable reconstruction and identification efficiency measurements. The upper boundary provides complementarity and exclusivity with respect to the hard-lepton and multi-lepton channels which facilitates the combination of the results. Events with subleading leptons are rejected to guarantee complementarity with the multi-lepton channels. The transverse mass requirement was fixed to $m_T > 100 \text{ GeV}$. While the transverse mass distribution provides substantial discriminating power between the Standard Model backgrounds and large- ΔM SUSY models, compressed

scenarios tend to be more background-like due to the soft leptons in the event. Attempts to increase the signal acceptance by lowering the m_T requirement or applying a two-sided m_T selection at both ends of the spectrum have not resulted in improved sensitivity. The transverse momentum requirement on subleading jets was set to the lowest experimentally feasible value of $p_T > 25$ GeV, since all jets in the event, with the exception of the hard initial state radiation jet, are expected to be soft.

The optimization with the variables described above was performed on one-step simplified model signal samples with gluino pair production and decays according to $\tilde{g}\tilde{g} \rightarrow qq\bar{q}\bar{q}WW\tilde{\chi}_1^0\tilde{\chi}_1^0$. The mass ratio parameter was set to $x = 1/4$. The samples were generated with HERWIG++ with 10 000 - 20 000 Monte-Carlo events per mass point. The non-optimal choice of generator and low available statistics resulted from delays in the Monte-Carlo production and the absence of better suited samples, such as the ones described in section 6.2.1 at the time when the optimization was performed. To determine the expected number of Standard Model background events, simulated samples such as the ones summarized in Tab. 6.2 were used. The simulated signal and background event numbers were scaled to an integrated luminosity of 5 fb^{-1} . The statistical significance for a discovery was determined for each model point taking into account the background prediction from simulation plus a flat 30 % uncertainty. For the calculation of the significance a Poisson approximated formula as suggested in [195] was used.

Fig. 6.15 shows the expected significance in the mass plane of the simplified models described above for two of the $3 \cdot E_T^{\text{miss}} \times 3 \cdot E_T^{\text{miss}}/m_{\text{eff}} \times 3 \cdot p_T^{\text{jet},1} \times 2 \cdot N_{\text{jet}} = 54$ cut combinations that have been tested for the electron (top) and muon channels (bottom). The black contour line encloses the area where a significance of three standard deviations is expected and the black numbers indicate the number of MC events available at a given mass point. It can be seen that the MC statistics are very limited. In particular in the compressed mass region close to the diagonal line it is difficult to draw meaningful conclusions.

While variations of the missing transverse energy, leading jet p_T , and jet multiplicity within the above cited intervals did not result into significant improvements of the sensitivity given the statistical MC uncertainties, a dependence on the $E_T^{\text{miss}}/m_{\text{eff}}$ selection could be observed. Fig. 6.15 (top) shows the discovery reach in the electron channel when selecting events with $E_T^{\text{miss}}/m_{\text{eff}} > 0.4$ (left) and $E_T^{\text{miss}}/m_{\text{eff}} > 0.3$ (right). While many features due to the limited MC statistics are visible in these figures it is apparent that the softer selection on $E_T^{\text{miss}}/m_{\text{eff}}$ increases the overall sensitivity in the compressed mass region. Likewise for the muon channel in the bottom part of Fig. 6.15, the discovery potential is improved with the softer $E_T^{\text{miss}}/m_{\text{eff}}$ requirement.

While a clear dependence on the jet multiplicity was not observed in the optimization procedure described above, $N_{\text{jet}} \geq 2$ was chosen to improve the signal acceptance. This choice is also supported by the fact that jets in compressed SUSY scenarios may be soft and fall below the experimentally viable transverse momentum threshold of 25 GeV used in this analysis. The missing transverse energy requirement was fixed to 250 GeV. This choice is sufficiently far away from the missing transverse energy trigger efficiency turn-on and allows for control and validation regions (see section 6.6) in between. For the leading jet- p_T threshold a value of 130 GeV has been chosen.

The chosen configuration of selection requirements corresponds to the right panels of Fig. 6.15.

Figure 6.15: Expected significance in the mass plane of a simplified model (see text) for the electron (top) and muon channels (bottom) for selection criteria as given in the figure. The black contour line encloses the area where a significance of three standard deviations is expected. The black numbers indicate the number of MC events available at a given mass point. The selection criteria applied to the figures on the right correspond to the ones chosen for the analysis.

6.5.3 Signal Region Definition

The final signal region selections of the soft-lepton analysis along with those of the complementary hard-lepton and multi-lepton channels are summarized in Tab. 6.8. For all analyses separate signal regions have been established for electrons and muons or combinations thereof in the multi-lepton channels. The selection criteria of the hard-lepton and multi-lepton channels will not be discussed further, since the corresponding search channels are not the main subject of this thesis. It is important to note, however, that all signal regions are mutually exclusive to facilitate a later combination of the results as will be shown section 6.10. The selection criteria of all signal regions are also illustrated in Fig. 6.16, where the cuts in the discriminating variables are shown in the $E_T^{\text{miss}} - H_T$ plane. H_T is defined as the scalar sum of the transverse momenta of the signal leptons, and jets in the event. Further details on the hard-lepton signal regions can be found in [166].

The performance of the soft-lepton signal region selection criteria was reevaluated in terms of acceptance $\times$ selection efficiency once the simulated signal samples with optimal generator setup and sufficient statistics as described in 6.2.1 became available. Figures 6.17 and 6.18 show the acceptance $\times$ selection efficiency values of the one-step simplified model for which the analysis

	soft-lepton	hard-lepton		multi-lepton	
	2-jet	3-jet	4-jet	2-jet	4-jet
Trigger	Missing E_T	Single electron or muon (+jet)		Single electron or muon (+jet)	
N_{lep}	1	1	1	≥ 2	≥ 2
E_T^{el} [GeV]	7 – 25	> 25	> 25	25	25
p_T^{μ} [GeV]	6 – 20	> 20	> 20	20	20
$E_T^{\text{el}_2}$ [GeV]	< 7	< 10	< 10	> 10	> 10
$p_T^{\mu_2}$ [GeV]	< 6	< 10	< 10	> 10	> 10
N_{jet}	≥ 2	≥ 3	≥ 4	≥ 2	≥ 4
p_T^{jet} [GeV]	$> 130, 25$	$> 100, 25, 25$	$> 80, 80, 80, 80$	$> 200, 200$	$> 50, 50, 50, 50$
$p_T^{\text{add, jet}}$ [GeV]	—	< 80	—	< 50	—
E_T^{miss} [GeV]	> 250	> 250	> 250	> 300	> 100
m_T [GeV]	> 100	> 100	> 100	—	—
$E_T^{\text{miss}}/m_{\text{eff}}$	> 0.3	> 0.3	> 0.2	—	0.2
$m_{\text{eff}}^{\text{inc}}$ [GeV]	—	> 1200	> 800	—	> 650

Table 6.8: Summary of the selection criteria of the exclusive signal regions of the soft-lepton, hard-lepton, and multi-lepton analyses [166].

is optimized as a function of the gluino-neutralino mass difference for chains of model points with same gluino mass. Fig. 6.17 (c)-(f) and Fig. 6.18 show that the acceptance $\times$ selection efficiency in the hard-lepton and multi-lepton signal regions decreases rapidly as the compressed region with small mass differences is approached. This explains the low sensitivity of previous analyses in this region as was shown for example in Fig. 6.1. In the soft electron and muon signal regions shown in Fig. 6.17 (a)-(b), however, the acceptance $\times$ selection efficiency remains high when approaching the compressed mass region.

Figure 6.16: Illustration of the selection criteria of three single-lepton signal regions (left) and the two multi-lepton signal regions in the $E_T^{\text{miss}} - H_T$ plane. H_T is defined as the scalar sum of the transverse momenta of the signal leptons, and jets in the event. [196]

Figure 6.17: Acceptance $\times$ efficiency in the soft-lepton (a)-(b) and hard-lepton (c)-(f) signal regions for a one-step simplified model as a function of the mass difference $m_{\tilde{g}} - m_{\tilde{\chi}_1^0}$ for chains of model points with same $m_{\tilde{g}}$. The acceptance $\times$ efficiency decreases rapidly in the hard-lepton signal regions as the compressed mass region is approached, while it remains high in the soft-lepton signal regions.

Figure 6.18: *Acceptance $\times$ efficiency in the multi-lepton signal regions for a one-step simplified model as a function of the mass difference $m_{\tilde{g}} - m_{\tilde{\chi}_1^0}$ for chains of model points with same $m_{\tilde{g}}$. The acceptance $\times$ efficiency decreases rapidly as the compressed mass region is approached.*

6.5.4 Cut-flows and Kinematic Distributions

Table 6.9 shows the number of events for data and simulated background samples after each of the preselection criteria described in section 6.5.1. The event numbers are combined for the electron and muon channels. A cut on the missing transverse energy consistent with the plateau region of the efficiency of the missing transverse energy triggers has been added in this table. The column “Other Bkg” contains the single-top, di-boson, $t\bar{t}$ + vector boson samples listed in Tab. 6.2. Agreement between data and simulated samples is not expected at these early stages of the selection due to the missing multijets component.

Tab. 6.10 shows the number of observed events and the expectation from simulated samples in events with one soft signal electron or muon according to tables 6.5 and 6.6. The event yields are shown after application of cuts on the jet multiplicity, $N_{\text{jets}} \geq 2$, and the transverse momentum threshold of the leading jet, $p_{\text{T}}^{\text{jet}_1} > 130$ GeV, on top of the selection criteria listed in Tab. 6.9. In this table the number of events in the electron and muon channels are listed separately. The column “Other Bkg” lists again the event numbers for the combined single-top, di-boson, and $t\bar{t}$ + vector boson samples. An additional column with a data-driven multijet estimate (to be described in section 6.6.4) has been added.

Figures 6.19 - 6.21 show comparisons of data and simulated samples after the jet multiplicity and leading jet p_{T} requirements shown in Tab. 6.10 for the most relevant kinematic distributions of the analysis. The yellow band includes systematic uncertainties on the jet energy scale and resolution, lepton identification and resolution, luminosity, multijet estimate, and, where appropriate, b-tagging. The exact definition of these uncertainties will be described in section 6.7. Within the given uncertainties, good agreement is observed between data and simulation in these distributions.

Cut	Data	W/Z+jets	Top	Other Bkg
No cut	49120355 $\pm$ 7008	57359020 $\pm$ 21083	280061 $\pm$ 191	199620 $\pm$ 146
GRL	49120355 $\pm$ 7009	57359020 $\pm$ 21083	280061 $\pm$ 191	199655 $\pm$ 146
Trigger	17994934 $\pm$ 4242	1267331 $\pm$ 1479	116680 $\pm$ 120	40638 $\pm$ 69
Jet cleaning	16904838 $\pm$ 4112	1264538 $\pm$ 1477	116468 $\pm$ 120	40557 $\pm$ 68
Smart veto	16459247 $\pm$ 4057	1263507 $\pm$ 1476	116250 $\pm$ 120	40506 $\pm$ 69
Primary Vertex	16294524 $\pm$ 4037	1259305 $\pm$ 1468	116136 $\pm$ 120	40380 $\pm$ 69
Bad/Cosmic muon veto	16288915 $\pm$ 4036	1258192 $\pm$ 1467	116064 $\pm$ 120	40336 $\pm$ 69
$E_{\text{T}}^{\text{miss}} > 180$ GeV	312730 $\pm$ 559	49589 $\pm$ 224	8980 $\pm$ 30	1712 $\pm$ 14

Table 6.9: Cut-flow table showing the number of events for data and the considered simulated samples as a function of the preselection criteria described in section 6.5.1 combining electron and muon channels. A cut on the missing transverse energy consistent with the plateau region of the efficiency of the missing transverse energy triggers has been applied. Agreement is not expected at this stage of the cut-flow due to the missing multijet component. [196]

Figure 6.19: Distribution of the lepton transverse momentum (top), the transverse mass (middle) and the missing transverse energy (bottom) in data and simulation including a data-driven multijet estimate for the electron channel (left) and muon channel (right) after the lepton, two jets, and $E_T^{\text{miss}} > 180$ GeV kinematic selection. The definition of the yellow uncertainty band is given in the text. [196]

Figure 6.20: Distribution of the number of jets (top), the transverse momentum of the leading (middle) and the second leading jet (bottom) in data and simulation including a data-driven multijet estimate for the electron channel (left) and muon channel (right) after the lepton, two jets, and $E_T^{\text{miss}} > 180$ GeV kinematic selection. The definition of the yellow uncertainty band is given in the text. [196]

Figure 6.21: Distribution of the effective mass (top) and the ratio $E_T^{\text{miss}}/m_{\text{eff}}$ (bottom) in data and simulation including a data-driven multijet estimate for the electron channel (left) and muon channel (right) after the lepton, two jets, and $E_T^{\text{miss}} > 180$ GeV kinematic selection. The definition of the yellow uncertainty band is given in the text. [196]

Channel	Cut	Data	Total Bkg	W/Z+jets	Top	Other Bkg	Multijet
Soft-electron	≥ 2 jets	2729 ± 52.2	2707.0 ± 36.7	2017.3 ± 33.5	542.0 ± 7.7	92.3 ± 3.4	55.4 ± 12.5
	$p_T^{\text{jet}_1} > 130$ GeV	2150 ± 46.4	2194.5 ± 33.5	1682.7 ± 30.7	385.4 ± 6.4	71.2 ± 3.1	55.2 ± 11.3
Soft-muon	≥ 2 jets	4114 ± 64.1	4365.9 ± 51.2	3255.8 ± 43.3	819.5 ± 9.8	143.6 ± 4.5	147.0 ± 25.2
	$p_T^{\text{jet}_1} > 130$ GeV	3241 ± 56.9	3491.3 ± 46.5	2667.8 ± 39.3	580.4 ± 8.2	111.5 ± 4.1	131.6 ± 23.0

Table 6.10: Cut-flow table showing the number of events for data and the considered simulated samples as well as a data-driven multijet estimate (see section 6.6.4). Event yields are shown for events with soft electrons or muons after the jet multiplicity and transverse momentum requirements on the leading jet, on top of the selection criteria applied in Tab. 6.9.

6.6 Background Estimation

This section describes the physical foundation of the background estimation methodology employed in this analysis. It is closely connected to section 6.7, which details the treatment of associated systematic uncertainties, and section 6.8, which describes how all components of the background estimation are combined into a global likelihood fit. It is important to emphasize that the backgrounds are determined for the soft-lepton, hard-lepton, and multi-lepton analyses at the same time using common control regions.

6.6.1 Overview

The dominant Standard Model backgrounds in the soft-lepton, as well as the hard-lepton, and multi-lepton analyses are events from W/Z +jets and $t\bar{t}$ processes. Backgrounds from multijet, di-boson, single-top, and $t\bar{t}$ + vector boson processes are present to a smaller extent. The technique used to estimate the dominant backgrounds is based on the enhancement of such events in dedicated control regions. Simulated samples are then normalized to the data in these control regions, using the jet multiplicity as the discriminating variable. The simulated samples are then used to extrapolate the prediction into the signal regions. The expected event yields are cross-checked in validation regions which lie between the control and signal regions.

The common control regions as well as the validation regions used in all three analyses are defined in section 6.6.3. The multijet background is estimated with an independent data-driven approach described in section 6.6.4. The remaining backgrounds are taken from simulation as pointed out in section 6.6.5. The fitting procedure is described in section 6.8.

Various discrepancies observed between data and simulation are addressed by a reweighting procedure for W +jets and Z +jets events which is performed prior to the fit in the W +jets, Z +jets, and $t\bar{t}$ control regions. The reweighting procedure is explained in section 6.6.2.

6.6.2 Reweighting of W +jets and Z +jets Simulated Samples

Discrepancies between data and simulation have been observed in the hard-lepton channels for a number of distributions, most notably missing transverse energy, effective mass, and jet transverse momentum. Fig. 6.23 (top left) shows the missing transverse energy distribution, where discrepancies of up to 40% and a significant overestimation of the simulated samples can be observed. In addition, the disagreement tends to get worse with increasing values of missing transverse energy.

In a preliminary version of the analysis [198] the discrepancies were found to be reduced if the W +jets simulated samples were reweighted at the level of the generated transverse momentum of the vector boson³. The reweighting factors were derived from a comparison of the generated vector boson p_T in ALPGEN (see Tab. 6.2) and SHERPA (v1.3.1) [197] generated samples. Fig. 6.22 shows a comparison between ALPGEN and SHERPA generated W +jets and Z +jets samples as a function of the true vector boson p_T . The ratio between the two generators indicates the magnitude of the correction that is applied to the ALPGEN samples. As can

³The transverse momenta of the vector bosons are defined as $p_T^W = \sqrt{(\vec{p}_{T,\text{lep1}} + \vec{p}_T^{\text{miss}})^2}$ and $p_T^Z = \sqrt{(\vec{p}_{T,\text{lep1}} + \vec{p}_{T,\text{lep2}})^2}$ for the W and Z bosons respectively.

Figure 6.22: Shape comparison of the generated W and Z boson transverse momentum spectra produced with ALPGEN (see Tab. 6.2) and SHERPA (v1.3.1) [197] W +jets and Z +jets samples. The middle panel shows that the ratio between the two predictions is very similar for the W +jets and Z +jets samples. The lower panel shows that similar shape differences between the W +jets and Z +jets samples are predicted by both generators. [196]

be seen, the correction factors for Z +jets simulated samples in bins of generated p_T^Z and for W +jets simulated samples in bins of generated p_T^W are very similar.

In the final version of the analysis, the generator-based reweighting procedure is replaced by a data-driven approach in which the correction factors are determined by fitting the reconstructed p_T^Z spectrum with the combined W +jets and Z +jets expectation from simulated samples binned in terms of generated vector boson p_T . The fitted normalizations of the combined W +jets and Z +jets samples are then taken as generated vector boson p_T dependent reweighting factors. The advantage of using only the reconstructed p_T^Z and not the p_T^W spectrum in this procedure is that p_T^Z is a purely leptonic quantity with little impact from systematic effects such as jet energy scale.

The fit is performed on an opposite-charge same-flavour di-lepton sample with invariant mass $80 \text{ GeV} < m_{ll} < 100 \text{ GeV}$ around the Z boson mass. The leptons are required to fulfill the signal lepton selection criteria listed in Tables 6.5 and 6.6. Furthermore, ≥ 3 signal jets with transverse momentum $p_T > 25 \text{ GeV}$ are required and a cut on the inclusive effective mass $m_{\text{eff}}^{\text{inc}} > 400 \text{ GeV}$ is applied. The combined W +jets and Z +jets simulated samples are binned in generated p_T^Z in intervals of 50 GeV between 0 and 200 GeV and one inclusive bin for values beyond 200 GeV. The normalizations of the small expected contributions from simulated $t\bar{t}$, single-top, and di-boson processes are fixed at their nominal cross-sections values. The multijet background contribution is determined with a data-driven method as described in section 6.6.4. Fig. 6.23 shows the reconstructed p_T^Z for the di-electron (top) and di-muon (bottom) selection before (left) and after the fit (right). As expected, the agreement between the distributions is significantly improved after the fit. The normalizations of the combined W +jets and Z +jets

Figure 6.23: Top: Distribution of E_T^{miss} in the hard-lepton W +jets control region for the combined electron and muons channels before (left) application of reweighting factors and after (right) the final fit to all background control regions [196]. Middle and Bottom: Reconstructed p_T^Z spectrum before (left) and after (right) the fit using simulated samples binned in generated vector boson p_T for di-electron (middle) and di-muon (bottom) event samples [166].

Generated $p_T^{W/Z}$ bin	Fitted reweighting factor
< 50 GeV	1.0 ± 0.0
50 GeV - 100 GeV	0.893 ± 0.043
100 GeV - 150 GeV	0.866 ± 0.056
150 GeV - 200 GeV	0.737 ± 0.061
> 200 GeV	0.772 ± 0.079

Table 6.11: *Fitted reweighting factors in bins of generated vector boson p_T .*

samples obtained from the fit are then used as generated vector boson p_T dependent reweighting factors for the simulated W +jets and Z +jets samples. Since only a shape correction and not the absolute normalization is of interest, the reweighting factors are normalized in the first bin of Z boson transverse momentum $0 \text{ GeV} < p_T^Z < 50 \text{ GeV}$. The final weights are shown in Tab. 6.11. They are commonly applied to W +jets and Z +jets simulated samples throughout the analysis.

Fig. 6.23 (top right) shows the missing transverse energy distribution after application of the reweighting factors and the final background fit described in section 6.8. The agreement between data and the reweighted simulated samples is significantly improved with respect to the initial comparison to simulation shown in Fig. 6.23 (top left).

6.6.3 Control and Validation Regions for $t\bar{t}$ and W/Z +jets Backgrounds

The definition of the control and validation regions used in the soft-lepton and hard-lepton (W +jets & $t\bar{t}$) as well as the multi-lepton (Z +jets & $t\bar{t}$) analyses are given in Table 6.12.

The soft-lepton control regions are defined by requirements on the transverse mass, $40 \text{ GeV} < m_T < 80 \text{ GeV}$, and the missing transverse energy, $180 \text{ GeV} < E_T^{\text{miss}} < 250 \text{ GeV}$. The lower missing transverse energy bound is governed by the efficiency turn-on of the E_T^{miss} -triggers (see Fig. 6.10). The requirements on the multiplicity and transverse momentum thresholds of jets and leptons are equivalent to those of the soft-lepton signal regions. The $t\bar{t}$ and W +jets control regions are then separated by the number of b-tagged jets found in the event. In the $t\bar{t}$ control region at least one of the two highest- p_T jets is required to be b-tagged. The W +jets control region is defined by events that did not fulfill this criterium. A combined W +jets and $t\bar{t}$ validation region is formed by combining the selection criteria of the two control regions and changing the missing transverse energy requirement to $80 \text{ GeV} < m_T < 100 \text{ GeV}$. For clarity, the position of the soft-lepton control, validation, and signal regions in the $m_T - E_T^{\text{miss}}$ plane is illustrated in Fig. 6.24 (top left). Fig. 6.25 (top and middle) shows a comparison of data and simulation including the data-driven multijet estimate (see section 6.6.4) for $E_T^{\text{miss}}/m_{\text{eff}}$ (top) and jet multiplicity (middle) in the soft-lepton W +jets (left) and $t\bar{t}$ (right) control regions.

The hard-lepton control regions are defined in a similar manner. The window in transverse mass is equivalent to that of the soft-lepton control regions, whereas for missing transverse energy lower values between 40 GeV and 150 GeV are chosen. This allows to collect more abundant statistics than in the soft-lepton control regions. In addition, events must contain at least three jets with transverse momentum thresholds of 80, 25, and 25 GeV and the inclusive effective mass is required to be greater than 500 GeV. The leptonic selections are equivalent

to the those in the hard-lepton signal regions. The distinction between the $t\bar{t}$ and W +jets control regions is again determined by the presence of b -tagged jets amongst the 3 hardest jets. The $t\bar{t}$ and W +jets validation regions are defined by shifting the missing transverse energy window to $150 \text{ GeV} < E_T^{\text{miss}} < 250 \text{ GeV}$. An additional high- m_T validation region is defined for $m_T > 100 \text{ GeV}$ inside a missing transverse energy window to $40 \text{ GeV} < E_T^{\text{miss}} < 250 \text{ GeV}$. The position of the hard-lepton control, validation, and signal regions in the $m_T - E_T^{\text{miss}}$ plane is illustrated in Fig. 6.24 (top right). As an example a comparison of data and simulation including the data-driven multijet estimate (see section 6.6.4) is shown in Fig. 6.25 (bottom left) as a function of jet multiplicity in the hard-lepton W +jets control region.

The multi-lepton control regions are formed by selecting events with at least two jets with transverse momentum thresholds of 80 and 50 GeV, or at least four jets with $p_T > 50 \text{ GeV}$. The leptonic requirements are the same as in the multi-lepton signal regions. The Z +jets

	soft-lepton		hard-lepton		multi-lepton	
	W CR	$t\bar{t}$ CR	W CR	$t\bar{t}$ CR	Z CR	$t\bar{t}$ CR
N_{jet}	≥ 2	≥ 2	≥ 3	≥ 3	≥ 2	≥ 2
p_T^{jet} [GeV]	$> 130, 25$	$> 130, 25$	$> 80, 25, 25$	$> 80, 25, 25$	$> 80, 50$ or $> 50, 50, 50, 50$	$> 80, 50$ or $> 50, 50, 50, 50$
$N_{\text{jet}} (b\text{-tagged})$	0	≥ 1	0	≥ 1	—	≥ 1
E_T^{miss} [GeV]	[180,250]	[180,250]	[40,150]	[40,150]	< 50	[30,80]
m_T [GeV]	[40,80]	[40,80]	[40,80]	[40,80]	—	—
$m_{\text{eff}}^{\text{inc}}$ [GeV]	—	—	> 500	> 500	—	—
$m_{\ell\ell}$ [GeV]	—	—	—	—	[81,101]	< 81 or > 101

	soft-lepton	hard-lepton		
	$W + t\bar{t}$ VR	W VR	$t\bar{t}$ VR	High- m_T VR
N_{jet}	≥ 2	≥ 3	≥ 3	≥ 3
p_T^{jet} (GeV)	$> 130, 25$	$> 80, 25, 25$	$> 80, 25, 25$	$> 80, 25, 25$
$N_{\text{jet}} (b\text{-tagged})$	—	0	≥ 1	—
E_T^{miss} (GeV)	[180,250]	[150,250]	[150,250]	[40,250]
m_T (GeV)	[80,100]	[40,80]	[40,80]	> 100
$m_{\text{eff}}^{\text{inc}}$ (GeV)	—	> 500	> 500	> 500

	multi-lepton 2-jet		multi-lepton 4-jet	
	Z VR	$t\bar{t}$ VR	Z VR	$t\bar{t}$ VR
N_{jet}	≥ 2	≥ 2	≥ 4	≥ 4
p_T^{jet} (GeV)	$> 120, 120$	$> 120, 120$	$> 80, 50, 50, 50$	$> 80, 50, 50, 50$
$N_{\text{jet}} (b\text{-tagged})$	—	≥ 1	—	≥ 1
E_T^{miss} (GeV)	[50,100]	[100,300]	[50,100]	[80,100]
$m_{\ell\ell}$ (GeV)	[81,101]	< 81 or > 101	[81,101]	< 81 or > 101

Table 6.12: *Top: Overview of the selection criteria for the W +jets, Z +jets, and $t\bar{t}$ control regions (CR). Only the selection criteria that are different from the signal selection criteria listed in Table 6.8 are shown. Middle and bottom: Overview of the selection criteria for the validation regions (VR). Only the selection criteria that are different from the control region selections are shown. [166]*

control region is defined by applying additional requirements on the missing transverse energy, $E_T^{\text{miss}} < 50$ GeV, and by selecting events inside the Z -mass window, $81 \text{ GeV} < m_{\ell\ell} < 101$ GeV. Events that do not fall into the invariant mass window, but fulfill $30 \text{ GeV} < E_T^{\text{miss}} < 50$ GeV and have at least one b -tagged jet amongst the above described jets, define the $t\bar{t}$ control region. Z +jets and $t\bar{t}$ validation regions are defined for the two-jet and four-jet selections. In the two jet selection the required jet p_T thresholds are 120 GeV each, in the four jet selection 80 GeV for the leading jet and 50 GeV for the 3 subleading jets. The usual b -jet requirements and invariant mass cuts are applied as for the control regions. The missing transverse energy window is $50 \text{ GeV} < E_T^{\text{miss}} < 100$ GeV for both Z +jets validation regions, $100 \text{ GeV} < E_T^{\text{miss}} < 300$ GeV for the two-jet $t\bar{t}$ validation region, and $80 \text{ GeV} < E_T^{\text{miss}} < 100$ GeV for the four-jet $t\bar{t}$ validation region. The position of the multi-lepton control, validation, and signal regions in the $m_T - E_T^{\text{miss}}$ plane is illustrated in Fig. 6.24 (bottom). As an example a comparison of data and simulation including the data-driven multijet estimate (see section 6.6.4) is shown in Fig. 6.25 (bottom right) as a function of jet multiplicity in the multi-lepton Z +jets control region.

Figure 6.24: Position of the soft-lepton (top left), hard-lepton (top right) and multi-lepton (bottom) control, validation, and signal regions in the $m_T - E_T^{\text{miss}}$ plane. Some of the validation regions in the multi-lepton channel have been suppressed. [196]

Figure 6.25: $E_T^{\text{miss}}/m_{\text{eff}}$ (top) and jet multiplicity (middle) distributions in the W +jets (left) and $t\bar{t}$ (right) control regions for data and simulation in the soft-lepton channel. Bottom: Jet multiplicity distributions for data and simulation in the hard-lepton W +jets (left) and multi-lepton Z +jets (right) control regions, respectively. For multijets a data-driven estimate is shown (see section 6.6.4).

6.6.4 Multijet Background Estimation

Due to the lepton requirement and the trigger-motivated high missing transverse energy selection, the multijet background in the soft-lepton analysis is small. It can arise from jets misidentified as leptons or from sufficiently isolated real leptons that are created as part of a jet. In addition, photons converted to electrons are counted as part of this background.

The three sources of backgrounds are estimated with a common data-driven approach which ATLAS-internally is referred to as “matrix-method”. The starting point of this method is the sample of preselected electrons and muons as defined in Tables 6.5 and 6.6. From this sample the number of leptons N_{pass} and N_{fail} is defined which pass or fail the signal selection criteria on top of the preselection requirements. N_{pass} and N_{fail} can be expressed in terms of the number of real (N_{real}) and misidentified ($N_{\text{misid.}}$) leptons in the preselected sample, according to:

$$N_{\text{pass}} = \epsilon_{\text{real}} N_{\text{real}} + \epsilon_{\text{misid.}} N_{\text{misid.}} \quad (6.5)$$

$$N_{\text{fail}} = (1 - \epsilon_{\text{real}}) N_{\text{real}} + (1 - \epsilon_{\text{misid.}}) N_{\text{misid.}} \quad (6.6)$$

The quantities ϵ_{real} and $\epsilon_{\text{misid.}}$ are defined as the efficiencies for real or misidentified leptons passing the preselection requirements to also pass the signal lepton criteria. The system of equations can be solved for the number of misidentified leptons in the signal lepton sample:

$$N_{\text{misid.}}^{\text{pass}} = \epsilon_{\text{misid.}} N_{\text{misid.}} = \frac{N_{\text{fail}} - (1/\epsilon_{\text{real}} - 1) N_{\text{pass}}}{1/\epsilon_{\text{misid.}} - 1/\epsilon_{\text{real}}} \quad (6.7)$$

It results from equation 6.7 that the multijet background can be obtained if the real lepton identification efficiency and the fake lepton misidentification efficiency are known.

The real lepton identification efficiency is measured for the full lepton p_T range with a T&P method in $Z \rightarrow ee$ and $Z \rightarrow \mu\mu$ events. The event sample is collected with single lepton (+jet) triggers as described in section 6.3. The events are required to contain 2 opposite-charge same-flavour leptons at preselection level (see Tables 6.5 and 6.6) with an invariant mass $80 \text{ GeV} < m_{\ell\ell} < 100 \text{ GeV}$ around the Z boson mass. W +jets events are rejected by requiring the missing transverse energy in the event to be lower than 30 GeV. At least one lepton, the so-called tag lepton, is required to fulfill the hard signal lepton selection criteria given in Tables 6.5 and 6.6 and must be matched to the single lepton trigger object in the event. Fig. 6.26 (top) shows the real lepton identification efficiencies ϵ_{real} used in the analysis as a function of lepton p_T , for electrons (left) and muons (right). The electron measurement is performed separately in bins of pseudorapidity since a small dependence on η was observed. The measurement in the calorimeter transition region $1.37 < |\eta| < 1.52$ is not used since electrons in this region are vetoed.

The misidentification efficiency is determined from a di-jet sample in which both jets are misreconstructed as leptons. The event sample is again collected with single lepton(+jet) triggers as described in section 6.3. Events with same-charge same-flavour lepton pairs at preselection level not falling into the Z boson mass window of $|m_{\ell\ell} - m_Z| > 25 \text{ GeV}$ are selected. At least one of the two leptons, the tag-lepton, is required to not fulfill the signal lepton requirements to enhance the fraction of di-jet events in the sample. The true lepton contributions, as predicted by simulation, are subtracted from the fake lepton sample. A 20 % systematic uncertainty on the cross-sections of the subtracted samples is propagated into $\epsilon_{\text{misid.}}$. Fig. 6.26 (bottom) shows

Figure 6.26: *Real lepton identification efficiencies (top) and misidentified lepton efficiencies (bottom) as a function of the lepton transverse momentum. The electron efficiencies are measured in bins of pseudorapidity. The measurement in the calorimeter transition region $1.37 < |\eta| < 1.52$ is not used since electrons in this region are vetoed.*

the efficiency for misidentified leptons, $\epsilon_{\text{misid.}}$, used in the analysis as function of lepton p_T , for electrons (left) and muons (right). The electron measurement has also been performed in bins of pseudorapidity.

The multijet estimate determined with the method described above has been applied in the distributions shown Figures 6.19 - 6.21. The good agreement between data and simulation demonstrated the validity of the estimates.

6.6.5 Other Backgrounds

Other minor backgrounds considered in the analysis are events from di-boson, single-top, and $t\bar{t}$ + vector boson production. Predictions of these backgrounds are taken from simulation using the MC samples and cross-sections listed in Tab. 6.2.

6.7 Treatment of Systematic Uncertainties

Systematic uncertainties affect the expected number of Standard Model background events as well as the signal predictions in the signal regions. The source of the uncertainties can

be of experimental nature, e.g. uncertainties on the determination of the energy scales of reconstructed objects, or inherent to the underlying theoretical model that describes the physics processes in the simulated samples. In this section the considered systematic uncertainties are described. They are then used as nuisance parameters of a given type in the combined background fit described in section 6.8.

6.7.1 Jet Energy Scale Uncertainty

A combination of several techniques, such as in-situ and test beam measurements of the calorimeter response to single particles and studies of simulated samples with systematic variations, are used in ATLAS to determine the uncertainties on the jet energy scale [103].

The jet energy scale uncertainty for the EM+JES calibration scheme used in this analysis is obtained in bins of jet transverse momentum and pseudorapidity from the `JetUncertaintyProvider` package [199] provided by the ATLAS jet / missing transverse energy performance group. Since events with rather large jet activity are expected in the analysis, additional contributions to the uncertainty from jet flavor composition and response, as well as close-by jet effects have to be taken into account. They are provided by the `MultijetJESUncertaintyProvider` package [200]. To determine the effect of the jet energy scale on the expected event yields in the signal regions, the simulated samples are processed with jet energies scaled up and down from the nominal values by the provided uncertainty.

The systematic jet energy scale uncertainty obtained by this procedure constitutes a large part of the yellow uncertainty band shown in Figures 6.19 - 6.21 and 6.25.

In the final background fit the jet energy scale uncertainties are grouped into three bins of jet transverse momenta at low (25-40 GeV), intermediate (40-100 GeV), and at high (> 100 GeV) values. The groups reflect jet- p_T regions where the components of the jet energy scale uncertainty contribute differently to the total uncertainty. At low values generator specific uncertainties related to hadronization and showering, effects from close-by jets, pile-up interactions, flavour composition and flavour response dominate, whereas at higher transverse momenta the calorimeter response uncertainties for single particles obtained from in-situ and test beam measurements gain importance. The uncertainties for each group are associated to a nuisance parameter in the combined fit (see section 6.8). Since the combined background fit is overconstrained, the uncertainties in the signal region can thus be expected to be constrained to some extent by the fit in the control regions. The degree to which a reduction of the input uncertainty is possible depends, however, on the transverse momentum spectrum of jets in the control and signal regions. If a given control region is populated only by low- p_T jets whereas the signal region contains a significant fraction of jets with higher transverse momenta, the effect will be limited.

6.7.2 Jet Energy Resolution

The jet energy resolution and its uncertainty are determined with in-situ techniques based on the transverse momentum conservation in dijet events [201]. The resulting numbers are encoded in the `JetEnergyResolutionProvider` tool [202], which provides the relative jet transverse momentum resolution σ_{p_T} and its uncertainty $\sigma_{\sigma_{p_T}}$ as a function of the jet transverse momentum

and rapidity. A smearing factor

$$\sigma_s = \sqrt{(\sigma_{p_T} + \sigma_{\sigma_{p_T}})^2 - \sigma_{p_T}^2} \quad (6.8)$$

is used to smear the transverse jet momentum by a Gaussian distribution with mean one and width σ_s . For each jet to be smeared, a random number is drawn from this distribution and multiplied with the original jet transverse momentum in the simulated samples. Due to the good agreement of the jet resolution between data and simulation no smearing of the central value is applied.

6.7.3 Lepton Uncertainties

As for jets, energy scale and resolution uncertainties are applied to leptons in simulated samples. The techniques used to derive these uncertainties are based on test beam measurements, measurements of the $Z \rightarrow \ell\ell$, $J/\psi \rightarrow \ell\ell$ resonances, and in-situ studies comparing the ID track momenta with the measured calorimeter energies (electrons) or track momenta in the MS (muons). These methods are described in detail in references [98] and [203].

The packages provided by the ATLAS electron / photon and muon combined performance groups are the **EnergyRescaler** [159] and **MuonMomentumCorrections** [204] tools. They are applied in a similar manner to that described in sections 6.7.1 and 6.7.2.

The scale factors correcting the lepton reconstruction and identification efficiencies described in sections 6.4.1 and 6.4.2 are also associated to systematic uncertainties of the corresponding efficiency measurements. The uncertainties are provided by the ATLAS electron / photon and muon Combined Performance Groups in references [205] and [204]. The simulated samples are processed with these uncertainties scaled up and down.

6.7.4 Missing Transverse Energy Uncertainty

The systematic energy scale and resolution uncertainties associated to jets and leptons as described above are propagated to the missing transverse energy. In addition, the systematic uncertainties associated to the “cell-out”-term of topological calorimeter clusters not belonging to any of the reconstructed objects in the E_T^{miss} calculation (see sections 4.3.4 and 6.4.4), are determined from simulated samples with systematic variations and from topocluster energy scale uncertainties as described in [104]. An additional uncertainty of 6.6 % of the total missing transverse energy is applied due to the impact of pileup interactions on the cell-out term. This uncertainty is extracted from the ratio of data and simulated samples in $Z \rightarrow \mu\mu$ events where no jets with transverse momentum greater than 20 GeV are reconstructed.

6.7.5 b-tagging Uncertainty

The scale factors correcting the b-tagging efficiency and the light flavour mis-tagging rate (see section 6.4.3) are subject to systematic uncertainties. The uncertainties associated to these measurements are provided by the ATLAS flavour tagging group [206] and are applied to the simulated samples with uncertainties scaled up and down respectively.

6.7.6 Trigger Uncertainty

The measurements of the electron and muon trigger efficiencies are subject to systematic uncertainties. These uncertainties are propagated to the trigger efficiency scale-factors that are used to bring data and simulation in agreement (see section 6.3).

6.7.7 Luminosity Uncertainty

The total integrated luminosity and its uncertainty is determined with van der Meer scans [207] and a number of luminosity-sensitive detectors (see section 3.2.4). The relative luminosity uncertainty has been determined to 3.9 % [208, 209].

6.7.8 Theoretical Background Uncertainties

The uncertainties related to the choice of the renormalization and factorization scales are usually evaluated by varying the corresponding scales up and down by a factor of two. These variations affect mostly the overall normalization of the cross-sections of the corresponding samples as a function of the number of generated partons in the process. In this analysis a different approach is used, in which the scale uncertainties are translated into a fit of the normalizations of the individual light parton bins of the ALPGEN generated W +jets, Z +jets, and $t\bar{t}$ samples. The details of the implementation are described in section 6.8.

Another generator related uncertainty arises from the MLM matching procedure described in section 4.4.1. The parton shower matching scale $p_T^{\min}$ determines down to which minimal transverse momentum value the partons are generated with the hard process matrix element calculation, and where the parton shower programs take over. The default choice for the ALPGEN generated samples is $p_T^{\min} = 15$ GeV. In line with the generator recommendations the uncertainty associated to this choice is evaluated by varying the scale up to $p_T^{\min} = 30$ GeV. The resulting difference is assigned as both a negative and positive uncertainty.

In a previous iteration of the analysis [164, 210], the impact on the Standard Model background expectation from initial and final state radiation effects was determined by studying $t\bar{t}$ simulated samples, generated with AcerMC [185] for the matrix element calculation, interfaced to PYTHIA [143] for parton showering. In these samples the initial and final state radiation is modelled by PYTHIA, where different tunes as suggested in [211] with varying amounts of radiation have been applied. From the observed variations of up to 10 % it was concluded that the initial state radiation should be modelled at matrix element level which is the case for the ALPGEN samples used in this analysis. The remaining systematic uncertainties are associated to the parton shower matching scale and are evaluated as described above.

The uncertainties related to the fragmentation and hadronization models are studied by comparing $t\bar{t}$ and W +jets samples produced with the POWHEG [144] generator for the matrix element calculation interfaced to PYTHIA or HERWIG for parton showering and fragmentation. Since the simulated $t\bar{t}$ and W +jets samples are used to extrapolate the expected number of events from the control to the signal regions, the observed differences between the two generator setups when going from control to signal region like selections are assigned as uncertainty. The effect is small in the soft-lepton analysis where the control regions lie close in E_T^{miss} to the signal regions. A deviation of the order of 2% for the $t\bar{t}$ and 1% for the W +jets samples is

found and applied as uncertainty. The effect is larger for the hard-lepton signal regions which are kinematically further separated from the corresponding control regions. The corresponding uncertainties can be found in [196].

Scaling factors and uncertainties on the production cross-sections of W and Z bosons in association with heavy flavour quarks are applied. Based on recent measurements [179], the cross-sections of the $Wb\bar{b}$ +jets and $Wc\bar{c}$ +jets samples listed in Table 6.2 are scaled up by 1.63 ± 0.76 and those of the $Wc\bar{c}$ +jets sample by 1.11 ± 0.35 . No correction factor is applied on the $Zb\bar{b}$ production cross-section but an uncertainty of $\pm 100\%$ is assigned.

6.7.9 Theoretical Signal Uncertainties

The procedure for the calculation of the signal cross-section uncertainties which arise from variations of the PDFs as well as the renormalization and factorization scales has been described in section 6.2.1.

Another important uncertainty on the signal samples in the context of this analysis is that associated to initial state radiation. As detailed in section 6.2.1, a first step towards a more realistic modeling of initial state radiation was the migration from HERWIG++ generated samples, using the parton shower approach, to the MADGRAPH+PYTHIA configuration, where one additional jet is generated at the matrix element level. The uncertainties are evaluated by generating additional samples with high and low radiation PYTHIA tunes and parton shower matching scales varied up and down by a factor of two from the nominal value. The event yields of these samples in the signal regions compared to those of the nominal samples are used to derive the uncertainty. Despite the large number of events generated for these samples, the low acceptances in the very compressed signal regions lead to significant statistical fluctuations in the calculated uncertainties. To reduce this effect the calculated uncertainties were compared when removing either the transverse mass or the missing transverse energy requirements from the signal region selections. In both cases similar uncertainties were found. The final uncertainty is calculated as the quadratic sum of the two values. The total error from initial state radiation was then determined by adding the maximum deviation from the different PYTHIA variations and parton shower matching scales in quadrature.

The procedure above described was performed for a number of points in the mass plane of the one-step simplified models described in section 6.2.1. As expected, an increase in the uncertainty was observed towards smaller mass differences between the initial squarks or gluinos and the LSP as well as at smaller gluino or squark masses. The relative uncertainty over the entire grid of mass points is thus parameterized as a function of the gluino or squark and LSP masses according to

$$\sigma_{ISR}(m_{\tilde{g}/\tilde{q}}, m_{\tilde{\chi}_1^0}) = 0.48 - 0.0007 \times m_{\tilde{g}/\tilde{q}} - 0.0003 \times (m_{\tilde{g}/\tilde{q}} - m_{\tilde{\chi}_1^0}). \quad (6.9)$$

All masses are given in units of GeV. If equation 6.9 becomes negative it is set to zero.

6.8 Combined Fit

All components of the background estimation (see section 6.6) and the associated systematic uncertainties (see section 6.7) are combined into a global fit based on the profile likelihood method. The likelihood function is built from binned probability density functions (PDF) in all regions of the fit. These PDFs include free parameters that adjust the relative contributions of the background and signal samples as well as constrained so-called “nuisance” parameters which are used to model the statistical and systematic uncertainties. The PDFs are fit to data to determine the optimal parameter values and errors. Assuming a sufficient number of measurements constraining the fit, not only the free parameters but also the nuisance parameters can be adjusted, which allows for a reduction of systematic uncertainties using data. In section 6.8.1, the general features of the utilized fitting framework are described. Section 6.8.2 then details the specific fit configurations used in the analysis.

6.8.1 Fitting Framework

The simultaneous fit performed in the analysis is implemented in the **HistFitter** package [212], which is based on **RooStats**’ [213] **Histfactory** software [214]. **Histfactory** is a tool to generate parameterized probability density functions in the **RooStats** framework based on **ROOT** histograms [96]. It allows for a flexible treatment of systematic uncertainties on the input histograms as will be demonstrated below. **HistFitter** is a wrapper tool around **HistFactory**. It provides added flexibility for the construction and analysis of PDFs and contains pre-configured tools for hypothesis testing and limit setting using **RooStats** functionality.

Following the notation used in [214], the general probability model constructed in **Histfactory** is given by

$$\mathcal{P}(\mathbf{n}, a_p | \phi_p, \alpha_p) = \prod_{c \in \text{channels}} \prod_{b \in \text{bins}} \text{Poisson}(n_{cb} | \nu_{cb}) \cdot \text{Gauss}(L_0 | \lambda, \Delta_L) \cdot \prod_{p \in \mathbb{S}} f_p(a_p | \alpha_p) \quad (6.10)$$

where the following index labelling scheme is used:

- c: Channel index labelling regions with disjoint selection criteria
- b: Bin index labelling the bins in a given channel
- p: Parameter index labelling the parameters of the fit
- s: Sample index labelling the samples used in the fit (appears below)

In addition two subsets of parameters are defined according to:

- $\mathbb{N} = \{\phi_p\}$, the unconstrained normalization parameters
- $\mathbb{S} = \{\alpha_p\}$, the “nuisance” parameters associated to external constraints

The term $\prod_{p \in \mathbb{S}} f_p(a_p|\alpha_p)$ in equation 6.10 is a constraint term describing auxiliary measurements a_p that constrain nuisance parameters α_p . The index p runs over the constrained parameters $\mathbb{S}$ which may or may not be channel or bin dependent. The vector $\mathbf{n}$ contains the number of observed events per channel and bin. The term $\text{Poisson}(n_{cb}|\nu_{cb})$ gives the Poisson probability to observe n_{cb} events in channel c and bin b if ν_{cb} events are expected. The expected number of events in a given channel and bin ν_{cb} is defined by the expression

$$\nu_{cb}(\phi_p, \alpha_p) = \sum_{s \in \text{samples}} \phi_{cs} \lambda_{cs} \gamma_{cb}(\boldsymbol{\alpha}) \eta_{cs}(\boldsymbol{\alpha}) \sigma_{csb}(\boldsymbol{\alpha}). \quad (6.11)$$

Here ϕ_{cs} is the product of unconstrained normalization factors for a given sample and channel defined as $\phi_{cs} = \prod_{p \in \mathbb{N}_c} \phi_p$. These parameters, or subsets thereof, are typically the parameters of interest in the fit, such as for example the signal and background sample normalizations.

λ_{cs} is a luminosity parameter for a given channel and sample, which can be set either to the nominal luminosity value L_0 , or can be a common parameter between all samples of a given channel which is constrained via a Gaussian penalty term $\text{Gauss}(L_0|\lambda, \Delta_L)$ with the luminosity uncertainty Δ_L .

γ_{cb} is a bin-dependent scale factor used for shape systematics (**ShapeSys**) and Monte-Carlo statistical uncertainties (**StatError**). The parameter is shared between the samples in a given channel and bin. **ShapeSys** is constrained independently for each bin with Gaussian penalty terms of given mean and width. For **StatError** a Poissonian distribution is assumed.

η_{cs} is a channel- and sample-dependent overall normalization factor (**OverallSys**). It is constrained with a Gaussian penalty term which contains the input uncertainty and two nuisance parameter for variations of mean and width.

σ_{csb} is a parameterized histogram for a given sample and channel (**HistoSys** and **normHistoSys**). The parameterization is a linear interpolation between the upward and downward scaled histograms corresponding to a given input uncertainty. The parameterization is thus fully correlated across the histogram bins. In the **HistoSys** configuration both shape and normalization of the histogram can be varied with one nuisance parameter each. Uncertainties can also be modelled as **normHistoSys**, where only the shape but not the normalization are allowed to vary.

To obtain the optimal values of the parameters of interest, e.g. the background normalizations in the signal regions, the logarithmic likelihood function derived from equation 6.10 for a given realization of the experimental data, is maximized.

6.8.2 Analysis Fit Setup

The specific likelihood function used in this analysis is constructed from the following previously discussed inputs:

- Observed number of events in the $t\bar{t}$ and W/Z +jets control regions, as well as the expectation from simulation in bins of jet multiplicity. In the soft-lepton analysis channel 6 bins with jet multiplicities ranging from 2 to 7 are used. In the hard-lepton and multi-lepton channels the jet multiplicities range from 3 to 9 and 2 to 9 respectively.

- Transfer factors from simulation to extrapolate the expected number of events in the control regions as obtained from the fit into the signal regions.
- Data-driven multijet estimate for the control and signal regions (see section 6.6.4).
- Background prediction from simulation for the remaining minor backgrounds from diboson, single-top, and $t\bar{t}$ + vector boson production.

The free parameters considered in the fit are the normalizations of the individual light parton bins (N_{parton}) of the ALPGEN $t\bar{t}$ and W/Z + jets samples. For the $t\bar{t}$ samples, each light parton bin from 0 to 3 is associated to one common free parameter for both semi-leptonic and di-lepton decays. The $t\bar{t}$ samples with four and five generated light partons share the same parameter with the $N_{\text{parton}} = 3$ sample. The W/Z +jets samples are associated to one common free parameter per light parton bin in the range $2 \leq N_{\text{parton}} \leq 5$. The normalizations of W/Z +jets samples with $N_{\text{parton}} \leq 2$ contribute only a low number of events due to the control and signal region selection criteria and are therefore fixed at their nominal values. The samples at $N_{\text{parton}} = 6$ share a common fit parameter with the $N_{\text{parton}} = 5$ samples. The W/Z plus heavy flavour samples share the same normalization parameter with the corresponding light flavour samples. The free parameters are labelled $\mu_x^{t\bar{t}}$ and μ_x^{WZ} , where x is the index for a given light parton bin.

The motivation behind fitting the light parton bin samples individually is the following: The dominant backgrounds $t\bar{t}$ and W/Z +jets are modelled with the ALPGEN generator at leading order. Higher order effects are usually estimated with variations of the renormalization and factorization scales as described in section 6.7.8. At leading order the uncertainties resulting from such variations are usually large. To reduce these uncertainties, all higher-order effects are absorbed into K-factors (i.e. normalizations) of the individual light parton bins that are determined from data. The fit is performed using the jet multiplicity distribution as discriminating variable, which is sensitive to the number of partons as can be seen in Fig. 6.27. For the method to work it is required that the shape of the jet multiplicity distribution in a given light parton bin is not distorted by scale variations or by changing from leading order to higher orders calculations. This has been verified with various generator configurations as further described in [196].

An additional free parameter in the fit, depending on the fit configuration (see below), is the normalization of the signal component, also referred to as signal strength. The corresponding parameter is labelled μ^{sig} .

The remaining background samples as well as the statistical and systematic uncertainties are constrained at some level and are thus modelled with nuisance parameters.

The first two light parton bins (0 and 1) of the W/Z +jets samples are, as mentioned above, not associated to free parameters due to the low statistics of these samples in the control and signal regions. Instead they are modelled with an `OverallSys` with Gaussian uncertainties determined from variations of the ALPGEN factorization and normalization scales varied up and down by a factor of 2. The corresponding nuisance parameters are labelled α_x^{WZ} , where x is again the index of a given light parton bin.

The uncertainty on the parton shower matching scale of the $t\bar{t}$ and W/Z +jets samples is modelled with a `normHistoSys` in all fit regions. As explained in section 6.7.8 the uncertainty is based on a variation of the matching scale from the nominal value of 15 GeV to 30 GeV. The

Figure 6.27: Jet multiplicity distributions for the light parton bins (0-5) of Alpen W +jets simulated samples [196].

observed difference in the prediction is also used to approximate the negative uncertainty. The corresponding nuisance parameters are labelled α_{mlm}^{WZ} and $\alpha_{\text{mlm}}^{t\bar{t}}$.

The hadronization uncertainty on the $t\bar{t}$ and W/Z +jets samples are modelled with one **OverallSys** each correlated in all regions. The nuisance parameters are labelled $\alpha_{\text{had}}^{t\bar{t}}$ and α_{had}^{WZ} .

The W/Z + heavy flavour samples are connected with an additional nuisance parameter α^{HF} which is modelled as **OverallSys** and is correlated over all light parton bins and reflects the uncertainty of the heavy flavour cross-section.

The backgrounds from single-top, di-boson, and $t\bar{t}$ + vector boson production and the associated uncertainties are commonly treated with a nuisance parameter $\alpha_{\text{err}}^{\text{BG}}$ implemented as an **OverallSys**. The normalization is set to constant such that only the uncertainty is allowed to vary.

The data-driven multijet estimate treated with an independent **HistoSys** for every fit region. The corresponding nuisance parameter is labelled $\alpha_{\text{r}}^{\text{multijet}}$, where r denotes the fit region. The normalization of the multijet background is fixed to the data-driven prediction in the validation and signal regions.

The initial state radiation uncertainties on the signal samples, which are only used in the exclusion fit configuration (see below), are modelled with an **OverallSys** in the relevant regions.

The jet energy scale uncertainty is modelled with 3 parameters as described in section 6.7.1. These parameters are of the type **normHistoSys** and are labelled $\alpha_{\text{low}}^{\text{JES}}$, $\alpha_{\text{medium}}^{\text{JES}}$, and $\alpha_{\text{high}}^{\text{JES}}$. If signal samples are considered in the fit separate jet energy scale uncertainty parameters are applied to these samples.

The lepton energy scale and resolution as well as the identification efficiency scale factor uncertainties are modelled with **OverallSys** parameters, which are correlated between channels of same lepton flavour. They are labelled α_c^{LE} (identification efficiency uncertainty), α_c^{LES} (energy scale uncertainty), and α_c^{LRI} (ID momentum resolution uncertainty), and α_c^{LRM} (MS momentum resolution uncertainty), where the index c denotes the channel or region.

The missing transverse energy cell-out and pile-up uncertainties are modelled with one **OverallSys** each for all samples and regions. The corresponding parameters are labelled $\alpha_{\text{cell-out}}^{\text{MET}}$ and $\alpha_{\text{pile-up}}^{\text{MET}}$.

The b-tagging uncertainty is modelled with one **OverallSys** parameter and is thus correlated in all relevant W/Z +jets and $t\bar{t}$ control region. The corresponding parameter is labelled $\alpha^{\text{b-tag}}$.

The uncertainties on the electron and muon trigger scale-factors are treated with **OverallSys** parameters which are correlated amongst regions that use the same trigger. The parameters are labelled $\alpha^{\text{trig,el}}$ and $\alpha^{\text{trig,mu}}$.

The uncertainties from the reweighting procedure described in section 6.6.2 are propagated into the fit. They are modelled by **OverallSys** parameters (labelled $\alpha_b^{\text{reweight}}$), where b stands for the bin index in true vector boson p_T .

The luminosity uncertainty is modelled as an **OverallSys** for the signal samples. The luminosity uncertainty on the combined minor backgrounds is absorbed in $\alpha_{\text{err}}^{\text{BG}}$. No assumptions on the luminosity are made for the other simulated samples used in the fit.

Finally also the statistical uncertainties from the simulated samples are taken into account. They are implemented with **StatError** parameters labelled $\gamma_{c,b}^{\text{stat}}$ in every region or channel c and jet multiplicity bin b . **HistFactory** does not assign this uncertainty for samples individually to avoid a multitude of nuisance parameters which could destabilize the fit.

The electron and jet energy resolution uncertainties discussed in section 6.7 are not propagated into the fit since their impact on the final result is found to be negligible and the presence of additional nuisance parameters associated to them destabilizes the fit.

With the above described setup several fit configurations are possible. In the so-called “background-only fit” only the control regions are used to constrain the fit parameters. Potential signal contributions are not taken into account. The expected backgrounds in the control regions are extrapolated to the signal regions using transfer factors from simulation. The validation regions, which lie between the control and signal regions, are used to test the adequacy of the background prediction.

In the “discovery fit” control regions and signal regions are used. The expected backgrounds in the control regions are again extrapolated to the signal regions using transfer factors from simulation. A possible signal contribution is considered in the signal region only. The background prediction in the signal region is conservative since potential signal contaminations of the control regions lead to an overestimation of the Standard Model backgrounds in the signal region. This fit configuration is used to derive model-independent upper limits on the visible cross-section of new physics scenarios.

In the “exclusion fit” configuration control and signal regions are used in the fit. The expected backgrounds in the control regions are again extrapolated to the signal regions using transfer factors from simulation. The signal contribution coming from a specific model prediction is fully taken into account in all regions. From this fit the exclusion limits for a given signal

model are derived. To improve the exclusion reach an additional shape fit in the signal region is performed. For this reason the requirement on the ratio of missing transverse energy and effective mass is removed and a fit is performed in this observable using the signal prediction.

6.9 Results

The background expectation in all regions of the analysis are determined with the background-only fit configuration. In this chapter the expected and observed event yields in the control, validation, and signal regions are summarized. In addition the systematic uncertainties in the signal region as derived from the fit are given.

6.9.1 Control Region Results

The background-only fit is performed simultaneously in all control regions. Thus all control regions have an influence on the Standard Model background predictions in the signal regions. The control regions with larger event samples typically have a higher weight in the fit. Fig. 6.28

Figure 6.28: W +jets control regions of the soft-lepton analysis for the electron (left) and muon channel (right) before (top) and after (bottom) application of the background-only fit.

shows the W control regions of the soft-lepton analysis for the electron (left) and muon channel (right) before (top) and after (bottom) the fit. As expected the agreement between the expected and observed number of events is improved after fit. Similarly Fig. 6.29 shows the top control regions of the soft-lepton analysis for the electron (left) and muon channels (right) before (top) and after (bottom) the fit. An improved agreement of the Standard Model expectation and the data is also visible here. The observed and expected numbers of events in all control regions are shown in Table 6.13 for the soft-lepton (top), hard-lepton (middle), and multi-lepton (bottom) control regions. The fitted number of events for each background component are listed separately. The quoted errors include the combined statistical and systematic uncertainties. In addition the raw number of events for each background component as used as input for the fit are shown. Fig. 6.30 gives an overview of the agreement between the fitted and the observed number of events for all control regions in terms of the pull distributions calculated as the difference between the observed and predicted number of events, divided by the combined statistical and systematic uncertainty.

Figure 6.29: $t\bar{t}$ control regions of the soft-lepton analysis for the electron (left) and muon channel (right) before (top) and after (bottom) application of the background-only fit.

Soft-lepton		Electron		Muon	
Number of events	$t\bar{t}$ CR	W CR	$t\bar{t}$ CR	W CR	
Observed	131	633	200		1092
Fitted bkg	131.0 ± 10.2	640.1 ± 20.9	194.4 ± 11.7		1093.5 ± 28.0
Fitted top	69.4 ± 5.8	31.8 ± 3.2	107.4 ± 6.4		57.7 ± 5.0
Fitted W/Z +jets	34.4 ± 4.4	593.9 ± 19.9	43.1 ± 6.2		998.7 ± 28.8
Fitted other bkg	10.6 ± 2.2	11.2 ± 2.3	16.2 ± 3.3		12.0 ± 3.2
Fitted multijet	16.5 ± 11.4	3.1 ± 0.5	27.7 ± 11.4		25.1 ± 21.1
MC exp. SM	125.7	679.3	212.8		1091.5
MC exp. top	78.8	44.3	122.6		72.5
MC exp. W/Z +jets	30.3	597.0	50.0		961.5
MC exp. other bkg	11.2	14.2	19.5		21.0
Data-driven multijet	5.4	23.9	20.7		36.4
Hard-lepton		Electron		Muon	
Number of events	$t\bar{t}$ CR	W CR	$t\bar{t}$ CR	W CR	
Observed	2225	4510	2137		4064
Fitted bkg	2270.9 ± 40.1	4515.3 ± 62.9	2130.8 ± 39.0		4060.4 ± 60.3
Fitted top	1842.2 ± 53.0	882.5 ± 63.7	1724.0 ± 50.5		806.8 ± 57.2
Fitted W/Z +jets	292.6 ± 40.7	3103.0 ± 89.3	255.0 ± 36.3		3116.3 ± 72.1
Fitted other bkg	119.6 ± 23.6	90.9 ± 19.0	101.2 ± 21.0		73.7 ± 15.8
Fitted multijet	16.6 ± 33.4	438.9 ± 107.6	50.6 ± 26.6		63.6 ± 68.2
MC exp. SM	2533.4	4932.1	2232.4		4317.3
MC exp. top	2012.4	1154.0	1816.7		1006.6
MC exp. W/Z +jets	315.5	3281.5	261.4		3119.8
MC exp. other bkg	132.2	115.7	116.4		104.1
Data-driven multijet	73.3	380.8	37.9		86.8
Multi-lepton		Electron-Electron		Electron-Muon	
Number of events	$t\bar{t}$ CR	Z CR	$t\bar{t}$ CR	Z CR	$t\bar{t}$ CR
Observed	574	8090	740	10915	1165
Fitted bkg	559.3 ± 16.1	8065.1 ± 91.8	706.9 ± 15.6	10937.5 ± 102.5	1168.0 ± 29.4
Fitted top	495.8 ± 13.2	67.8 ± 1.9	648.3 ± 15.9	82.2 ± 2.2	1129.5 ± 31.2
Fitted W/Z +jets	36.5 ± 0.9	7924.2 ± 104.2	40.5 ± 1.0	10763.6 ± 102.8	3.2 ± 0.3
Fitted other bkg	11.8 ± 2.3	69.8 ± 14.0	16.6 ± 3.3	91.2 ± 18.3	26.4 ± 5.3
Fitted multijet	15.1 ± 14.9	3.2 ± 37.4	1.6 ± 2.7	0.5 ± 4.0	8.9 ± 13.7
MC exp. SM	510.1	7393.4	678.6	10273.0	1141.5
MC exp. top	461.3	65.1	619.4	80.9	1101.5
MC exp. W/Z +jets	34.1	7253.2	40.9	10093.8	3.4
MC exp. other bkg	11.5	73.0	17.0	98.0	28.5
Data-driven multijet	3.2	2.2	1.2	0.3	8.1

Table 6.13: Observed and expected event yields in the soft-lepton (top), hard-lepton (middle), and multi-lepton (bottom) control regions. The fitted number of events are given with combined statistical and systematic uncertainties. In addition the inputs to the fit from simulation as well as the data-driven multijet estimate are shown.

Figure 6.30: Summary of the fit results in the control regions. The difference between the observed and expected number of events, divided by the combined statistical and systematic uncertainty on the expectation, is shown for the soft-lepton, hard-lepton, and multi-lepton analyses.

6.9.2 Validation Region Results

To validate the background estimation procedure, the background estimates obtained in the control regions are extrapolated to so-called validation regions before unblinding the expected and observed event yields in the signal regions. The validation regions lie between the control and signal regions as defined in section 6.6.3. Fig. 6.31 shows combined $W + t\bar{t}$ validation regions for the electron (left) and muon channel (right) before (top) and after (bottom) the fit. The corresponding observed and predicted event yields in the soft-lepton validation regions are shown in Tab. 6.14 along with their combined statistical and systematic uncertainties. For information

Soft-lepton	Electron	Muon
Number of events	$W + t\bar{t}$ VR	$W + t\bar{t}$ VR
Observed events	295	463
Fitted bkg events	329.0 ± 17.6	478.1 ± 22.9
Fitted top events	60.2 ± 3.6	71.2 ± 3.3
Fitted WZ events	245.1 ± 10.5	374.5 ± 13.7
Fitted other BGs events	10.3 ± 2.1	14.9 ± 3.0
Fitted multijet	13.4 ± 11.7	17.5 ± 18.2
MC exp. SM events	342.4	479.1
MC exp. top events	69.5	80.9
MC exp. WZ events	247.1	363.7
MC exp. other BGs events	12.4	17.1
Data-driven multijet	13.4	17.5

Table 6.14: Observed and expected event yields in the soft-lepton validation regions. The fitted number of events are given with the combined statistical and systematic uncertainties. In addition the inputs to the fit from simulation as well as the data-driven multijet estimate are shown.

Figure 6.31: Validation regions of the soft-lepton analysis for the electron (left) and muon channel (right) before (top) and after (bottom) application of the background-only fit.

Figure 6.32: Summary of the fit results in the validation regions. The difference between the observed and expected number of events, divided by the combined statistical and systematic uncertainty on the expectation, is shown for the soft-lepton, hard-lepton, and multi-lepton analyses.

the raw event numbers as used as input for the fit are shown as well. The expected number of events is slightly higher than the observation in the electron channel, whereas consistency within one standard deviation is observed in the muon channel. Fig. 6.32 shows the pull distributions calculated as the difference between the observed and predicted number of events, divided by the combined statistical and systematic uncertainty for all validation regions used in all analyses, including the hard-lepton and multi-lepton channels. The corresponding observed and expected event yields are not explicitly relevant for this analysis and are therefore listed in Appendix B.1.

6.9.3 Signal Region Results

Fig. 6.33 shows the observed and expected number of events in the soft electron (left) and soft muon channel (right) before (top) and after (bottom) the fit as a function of the ratio of missing transverse energy and effective mass, where the signal region selection requirement on that observable has been removed for illustrative purposes. The corresponding observed and expected event yields in the signal regions are listed in Tab. 6.15 along with their combined statistical and systematic uncertainties. For information the raw event numbers as used as input for the fit are also shown. The observed and expected event yields of the hard-lepton and multi-lepton signal regions are not explicitly relevant for this analysis and are therefore listed in Appendix B.2. No significant excess is observed in any of the soft-lepton, hard-lepton, or multi-lepton signal regions.

Soft-lepton	Electron	Muon
Number of events	Signal region	Signal region
Observed	11	14
Fitted bkg	14.0 ± 3.3	19 ± 5
Fitted top	3.8 ± 0.6	3.8 ± 0.8
Fitted W/Z +jets	5.8 ± 1.0	11.4 ± 2.3
Fitted other bkg	0.6 ± 0.1	0.2 ± 0.1
Fitted multijet	3.8 ± 2.5	3.6 ± 2.5
MC exp. SM	14.2	18.0
MC exp. top	4.3	3.8
MC exp. W/Z +jets	5.5	10.5
MC exp. other bkg	0.5	0.1
Data-driven multijet	3.8	3.6

Table 6.15: *Observed and expected event yields in the soft-lepton signal regions. The fitted number of events are given with the combined statistical and systematic uncertainties. In addition the inputs to the fit from simulation as well as the data-driven multijet estimate are shown.*

Figure 6.33: *Distribution of the number of jets in the soft-lepton signal region in the electron (left) and muon (right) final state after applying the background-only fit.*

6.9.4 Systematic Uncertainties

The systematic uncertainties affecting the event yields in the signal regions are evaluated individually for each parameter. This is done by setting all parameters except the parameter of interest to constant in the fit and calculating the error propagated due to that parameter. Tab. 6.16 shows the full list of systematic uncertainties in units of expected events. The total background uncertainty in the signal region amounts to approximately 24% in the electron and muon channels. The largest contributions are statistical uncertainties (approximately 15-20%), uncertainties related to the multijet background estimation (approximately 13-18%), and uncertainties from the reweighting procedure of the W/Z +jets and $t\bar{t}$ simulated samples (approximately 4 %). The experimental uncertainties such as the jet energy scale uncertainties, missing transverse energy uncertainties, and uncertainties in the lepton identification efficiency correction factors contribute at a lower level.

Soft-lepton Number of Events	Electron Signal region	Muon Signal region
Total background uncertainty	± 3.28	± 4.47
$\mu_{\text{Np0}}^{t\bar{t}}$	± 0.02	± 0.00
$\mu_{\text{Np1}}^{t\bar{t}}$	± 0.04	± 0.04
$\mu_{\text{Np2}}^{t\bar{t}}$	± 0.12	± 0.10
$\mu_{\text{Np3}}^{t\bar{t}}$	± 0.15	± 0.19
μ_{Np2}^{WZ}	± 0.16	± 0.53
μ_{Np3}^{WZ}	± 0.14	± 0.11
μ_{Np4}^{WZ}	± 0.09	± 0.09
μ_{Np5}^{WZ}	± 0.00	± 0.15
α_{Np0}^{WZ}	± 0.00	± 0.00
α_{Np1}^{WZ}	± 0.00	± 0.00
$\alpha_{\text{soft-electron}}^{\text{multijet}}$	± 2.45	± 0.00
$\alpha_{\text{soft-muon}}^{\text{multijet}}$	± 0.00	± 2.43
$\alpha_{\text{err}}^{\text{BG}}$	± 0.11	± 0.05
$\alpha_{\text{had}}^{t\bar{t}}$	± 0.06	± 0.06
α_{had}^{WZ}	± 0.05	± 0.09
$\alpha_{\text{mlm}}^{t\bar{t}}$	± 0.00	± 0.00
α_{mlm}^{WZ}	± 0.00	± 0.00
$\alpha_{\text{JES}}^{\text{low}}$	± 0.13	± 0.14
$\alpha_{\text{JES}}^{\text{medium}}$	± 0.12	± 0.02
$\alpha_{\text{JES}}^{\text{high}}$	± 0.02	± 0.01
$\alpha_{\text{b-tag}}$	± 0.00	± 0.00
$\alpha_{\text{soft-electron}}^{\text{LES}}$	± 0.00	± 0.00
$\alpha_{\text{soft-muon}}^{\text{LES}}$	± 0.00	± 0.00
$\alpha_{\text{soft-electron}}^{\text{LE}}$	± 0.32	± 0.00
$\alpha_{\text{soft-muon}}^{\text{LE}}$	± 0.00	± 0.01
$\alpha_{\text{soft-muon}}^{\text{LRI}}$	± 0.00	± 0.05
$\alpha_{\text{soft-muon}}^{\text{LRM}}$	± 0.00	± 0.00
$\alpha_{\text{MET}}^{\text{cell-out}}$	± 0.01	± 0.07
$\alpha_{\text{pile-up}}^{\text{MET}}$	± 0.22	± 0.01
$\alpha_{\text{trig,el}}$	± 0.00	± 0.00
$\alpha_{\text{trig,mu}}$	± 0.00	± 0.00
α_{HF}	± 0.10	± 0.36
$\alpha_{\text{reweight}}^{200\text{GeV}}$	± 0.00	± 0.04
$\alpha_{\text{reweight}}^{150\text{GeV}}$	± 0.00	± 0.02
$\alpha_{\text{reweight}}^{100\text{GeV}}$	± 0.42	± 0.73
$\alpha_{\text{reweight}}^{50\text{GeV}}$	± 0.00	± 0.00
$\gamma_{\text{stat}}^{\text{soft-electron}}$	± 2.14	± 0.00
$\gamma_{\text{stat}}^{\text{soft-muon}}$	± 0.00	± 3.72

Table 6.16: Breakdown of the uncertainties on the expected background event yields in the soft-lepton signal regions in electron (left) and muon (right) channels obtained from the background-only fit configuration. The individual uncertainties can be correlated and do not necessarily add up quadratically to the total background uncertainty.

6.10 Interpretation

In this section exclusion limits are derived using the CL_s method [215]. Section 6.10.1 describes the statistical procedure used to compute these exclusion limits. In section 6.10.2 the model-independent upper limits on the visible cross-section for beyond the Standard Model events are given. In section 6.10.3 the derivation of exclusion limits for the specific supersymmetric models targeted with this analysis is described.

6.10.1 Statistical Procedure

For the statistical interpretation of experimental results the p-value is used to quantify the level of agreement of the observed data with a given hypothesis. In general, the p-value is defined as the probability to find data at least as incompatible with a given hypothesis as the observed data assuming that the hypothesis is true. Following the frequentist approach used at the LHC [216], the p-value can be determined by defining a test statistic

$$t_\mu = -2\ln\lambda(\mu), \quad (6.12)$$

where $\lambda(\mu)$ is the likelihood ratio defined as

$$\lambda(\mu) = \frac{L(\mu, \hat{\boldsymbol{\theta}})}{L(\hat{\mu}, \hat{\boldsymbol{\theta}})}. \quad (6.13)$$

$L(\mu, \boldsymbol{\theta})$ is the likelihood function of the fit as described in section 6.8, where μ is the free signal strength parameter and $\boldsymbol{\theta}$ is the vector of remaining free and constrained parameters of the fit. The background-only hypothesis corresponds to a value $\mu = 0$ whereas the signal-plus-background hypothesis is represented by $\mu = 1$ for a given signal model. $\hat{\boldsymbol{\theta}}$ in the numerator of equation 6.13 corresponds to the values of the fit parameters that maximizes the likelihood for a given value of the signal strength μ . $\hat{\mu}$ and $\hat{\boldsymbol{\theta}}$ in the denominator are the parameter values that maximize the likelihood if the signal strength is left as a free parameter in the fit. The likelihood ratio can take values $0 \leq \lambda \leq 1$. If the value of λ is close to one the observed value of the signal strength $\hat{\mu}$ is close to the hypothesized value of μ . The test statistics t_μ is thus a measure of the disagreement between the hypothesis and the data from which the p-value for that hypothesis can be calculated according to:

$$p_\mu = \int_{t_{\mu, obs}}^{\infty} f(t_\mu|\mu) dt_\mu \quad (6.14)$$

Here $t_{\mu, obs}$ is the value of the test statistic in data and $f(t_\mu|\mu)$ is the probability density function of the test statistic for a given value of μ .

For a particular signal model one can now determine the probability under the assumption of the signal-plus-background hypothesis ($\mu = 1$) to observe a value of the test statistic at least

as extreme as the one actually observed:

$$p_1 = p_{s+b} = \int_{t_{\mu,obs}}^{\infty} f(t_{\mu}|s+b) dt_{\mu} \quad (6.15)$$

Equally, one can determine the probability under the assumption of the background-only hypothesis ($\mu = 0$), to find values of the test statistic that are at least as incompatible as the observed one:

$$p_0 = p_b = \int_{t_{-\infty}}^{t_{\mu,obs}} f(t_{\mu}|b) dt_{\mu} \quad (6.16)$$

This is illustrated in Fig. 6.34 (left) which shows the value of a test statistic distribution for the signal-plus-background and the background-only hypotheses, as well as the observed value. The p-values for the signal-plus-background and background-only hypotheses correspond to the red and green shaded areas.

For one-sided (upper) exclusion limits the test statistic of equation 6.12 is set to zero for values $\hat{\mu} > \mu$ since upward fluctuations of the data are not considered as incompatibility with the assumed value of μ . In the exclusion case the test statistic t_{μ} in equations 6.15 - 6.16 is thus replaced by:

$$q_{\mu} = \begin{cases} -2 \ln \lambda(\mu) & | \hat{\mu} \leq \mu, \\ 0 & | \hat{\mu} > \mu. \end{cases} \quad (6.17)$$

Figure 6.34: Illustration of the test statistic distributions for the background-only and signal-plus-background hypotheses for the case of higher (left) and lower (right) signal expectations resulting in more or less separated distributions. A downward fluctuation of the data (right) can lead to a spurious exclusion if the CL_s method is not used. In this particular example the test statistic is defined such that it can become negative.

To exclude a particular signal model at 95 % confidence level it is typically required that the p-value associated to the signal-plus-background hypothesis satisfies the relation

$$p_{s+b} < \alpha, \quad (6.18)$$

where $\alpha = 1 - 0.95 = 0.05$. This approach is called the CL_{s+b} method.

However, if the number of signal events of a given model is very small and the test statistic of the signal-plus-background and the background-only hypotheses are similar as shown in Fig. 6.34 (right), a downward fluctuation of the observed event yield (corresponding to an upward fluctuation of the observed value of the test statistic in Fig. 6.34) can lead to the false exclusion of models to which the underlying analysis is not sensitive.

To avoid this undesired effect the CL_{s+b} method is replaced by the conservative CL_s approach [215], where a CL_s value is calculated according to

$$\text{CL}_s = \frac{p_{s+b}}{1 - p_b} < \alpha. \quad (6.19)$$

If the distributions of the test statistics of the signal-plus-background and background-only hypotheses are well separated, the value of $(1 - p_b)$ in the denominator of equation 6.19, also referred to as CL_b , is close to one and has only little impact on the exclusion power. If the two distributions lie closely together, and there is little sensitivity to the signal model, the numerator and the denominator of the equation 6.19 have similar values, and the CL_s value approaches unity, resulting in the non-exclusion of the signal model. Due to the low signal acceptance and the background-like distributions of compressed supersymmetric models, the usage of the CL_s method for the determination of exclusion limits is of particular relevance for the interpretation of the results of this analysis.

The test statistic q_μ is typically evaluated with pseudo-experiments. It is also possible to determine the p-values needed for limit setting with fast asymptotic formulae. In the asymptotic approach it follows from results by Wald [217] and Wilks [218] that in the large sample limit the test statistic in equation 6.17 can be approximated by a χ^2 distribution with one degree of freedom and then reads:

$$q_\mu = \begin{cases} \frac{(\mu - \hat{\mu})^2}{\sigma^2} & | \hat{\mu} \leq \mu, \\ 0 & | \hat{\mu} > \mu. \end{cases} \quad (6.20)$$

It can be shown that in this case (see [216] and references therein) the p-value for a given hypothesized value of μ is given by

$$p_\mu = 1 - \Phi(\sqrt{q_\mu}), \quad (6.21)$$

where Φ is the cumulative distribution of a normal distribution with zero mean and unit variance.

6.10.2 Upper Limits on Visible Signal Cross-section

Given the absence of a significant excess of events in any of the signal regions, the discovery fit configuration as described in section 6.8.2 is used to place model-independent upper limits on

Signal channel	S_{obs}^{95}	S_{exp}^{95}	$\langle\epsilon\sigma\rangle_{\text{obs}}^{95}[\text{fb}]$	CL_b
soft electron	$10.4^{+4.2}_{-3.1}$	8.6	1.82	0.28
soft muon	$12.5^{+5.4}_{-3.8}$	9.0	1.92	0.21

Table 6.17: Observed 95% CL upper limits on the visible cross section ($\langle\epsilon\sigma\rangle_{\text{obs}}^{95}$) as well as observed (S_{obs}^{95}) and expected (S_{exp}^{95}) upper limits on the number of signal events in the soft-lepton signal regions. In addition the CL_b value, i.e. the observed confidence level for the background-only hypothesis, is shown.

the visible signal cross-section. For this purpose a generic signal strength parameter μ in units of beyond Standard Model events is introduced in the signal regions. A scan of the corresponding CL_s values as a function of a given value of μ is performed to determine the lowest number of beyond Standard Model events that can be excluded at 95 % confidence level. The results are summarized in Table 6.17. The left two columns show the observed (S_{obs}^{95}) and expected (S_{exp}^{95}) 95 % confidence level limits on the number of beyond Standard Model events in the signal regions. By dividing S_{obs}^{95} by the integrated luminosity of the data one obtains a limit $\langle\epsilon\sigma\rangle_{\text{obs}}^{95}$ on the visible cross-section for beyond Standard Model events in the signal regions. These numbers are conservative since signal contaminations in the control regions which are not taken into account would lead to an overestimation of the backgrounds in the signal regions. The CL_b values, i.e. the confidence levels for the background-only hypothesis as defined in the previous section, are also listed in Table 6.17 and demonstrate good compatibility with the hypothesis. The model-independent upper limits on the visible signal cross-section of the hard-lepton and multi-lepton analyses are listed in Appendix B.3.

6.10.3 Exclusion limits

Exclusion limits are derived using the exclusion fit configuration described in section 6.8.2 for the simplified models described in section 6.2.1. To calculate the limit for a given model at its nominal cross-section, the CL_s value is evaluated for signal strength $\mu = 1$. To determine the excluded cross-section for a given model, 20 values of signal strength are scanned inside a region of interest where the CL_s value is expected to cross the threshold of 0.05. The lowest excluded signal strength is then taken as the point where the interpolated CL_s value reaches the 0.05 threshold. This is illustrated in Fig. 6.35 which shows the expected and observed CL_s values as a function of the signal strength for a one-step simplified model with gluino production and large mass difference between gluino and the LSP (left) as well as a model with a compressed mass spectrum (right). As expected the sensitivity is reduced in the compressed scenario. It can also be seen from Fig. 6.35 that the CL_s procedure is more conservative than the CL_{s+b} procedure, in particular at low signal strengths as previously discussed in section 6.10.1.

To increase the sensitivity, the exclusion fit does not only take into account the signal contamination in the control regions, but also performs an additional shape fit in the signal region. The discriminating variable for this fit is the ratio of the missing transverse energy over the effective mass, where the requirement on that observable has been removed from the signal-region selection shown in Table 6.8. The fit is performed in 6 bins of $E_T^{\text{miss}}/m_{\text{eff}}$ between 0.1 and 0.7 as shown in Fig. 6.33.

Figure 6.35: Observed and expected CL_s , CL_b , and CL_{s+b} values scanned as a function of the assumed signal strength μ for two one-step simplified models with gluino pair production and mass ratio parameter $x = 1/2$. A scenario with large mass difference between the gluino and neutralino is shown on the left and a model with a compressed mass spectrum on right.

The results for the simplified models described in section 6.2.1 are shown in figures 6.36 - 6.40 as a function of the parameters of the models. The solid red lines correspond to the observed exclusion contours assuming 100 % branching fraction of the pair produced initial particles into the decays of the given simplified models. The dotted red lines correspond to the $\pm 1\sigma$ variations of the cross-sections of the SUSY models calculated according to the procedure described in 6.2.1. To obtain the uncertainty band the exclusion fit is repeated three times, once for the nominal cross-section value and once with the theoretical uncertainties varied up and down respectively. The interlaced line around the yellow band corresponds to the expected limit and its $\pm 1\sigma$ variation. The uncertainty band also includes the initial state radiation uncertainty on the signal samples, as described in section 6.7.9. The grey numbers in the plots give the excluded cross-section in units of fb. All limits are determined with the CL_s procedure at 95 % confidence level. The sensitivity inside the parameter plane of the considered scenarios is governed by their production cross-sections as well as the product of acceptance times selection efficiency in the signal regions. For reference these numbers can be found in Appendix B.4.

Fig. 6.36 shows the exclusion contours and excluded cross-section limits for the one-step simplified models with gluino pair production and mass ratio parameter $x = 1/2$. As described in section 6.5.2 the soft-lepton signal regions have been specifically optimized for the compressed region of this scenario. In the top left panel the exclusion limits are shown for the combined soft-lepton signal regions, in the top right for the combined hard-lepton signal regions, in the bottom left for the combined multi-lepton signal regions, and in the bottom right for the combination of all signal regions. It can be seen that the soft-lepton analysis adds significant sensitivity in the compressed mass region close to the diagonal line where the gluino and LSP masses lie closely together. The observed limit along the compressed region now extends up to approximately 550 GeV in gluino mass, and the sensitivity in terms of excluded cross-sections is improved up to a factor of 10 throughout the entire compressed mass region when compared to the hard-lepton and multi-lepton channels.

Fig. 6.37 shows the exclusion contours and excluded cross-section limits for the one-step simplified models with gluino pair production and variable x parameter. Again, in the top left panel

the exclusion limits are shown for the combined soft-lepton signal regions, in the top right for the combined hard-lepton signal regions, in the bottom left for the combined multi-lepton signal regions, and in the bottom right for the combination of all signal regions. The soft-lepton signal regions are sensitive up to gluino masses of approximately 600 GeV. Naively a larger sensitivity at low values of x would be expected, due to the small chargino-neutralino mass differences resulting in soft-leptons from the W decay. However, small values of x also imply larger gluino-chargino mass differences which can result in a transverse boost of the chargino and eventually of the lepton in the subsequent decay. Thus the sensitivity for this simplified model is expected to be better in the hard-lepton signal regions.

Fig. 6.38 shows the exclusion contours and excluded cross section limits for the one-step simplified models with squark pair production and $x = 1/2$ (left) as well as variable x parameter (right). The limits are less stringent than in the case of gluino pair-production due to the lower production cross-sections. The combined exclusion limits for all signal regions are shown.

The exclusion contours and excluded cross-section limits for the two-step simplified models with gluino and squark pair-production are shown in Fig. 6.39 and Fig. 6.40 respectively. The topology of these scenarios is described in section 6.2.1. The soft-lepton signal regions have not been explicitly optimized for these scenarios, however, they were found to add some sensitivity in the compressed mass regions. This can also be inferred from the acceptance $\times$ selection efficiency numbers in Appendix B.4.

Figure 6.36: *Exclusion contours and excluded cross-section limits at 95% confidence level in the parameter space of one-step simplified models with gluino pair production and decays according to $\tilde{g} \rightarrow q\bar{q}\tilde{\chi}_1^\pm \rightarrow q\bar{q}'W^\pm\tilde{\chi}_1^0$. The chargino mass is fixed half way between the gluino and neutralino masses ($x=1/2$). The limits are shown for the soft-lepton (top left), hard-lepton (top right), and multi-lepton signal regions (bottom left), as well as for the combination of all signal regions (bottom right). The band around the median expected limit shows the $\pm 1\sigma$ variations, including all uncertainties except theoretical uncertainties on the signal. The dotted lines around the observed limit correspond to the $\pm 1\sigma$ variations of the theoretical uncertainties of the signal samples. The grey numbers indicate the excluded cross-section in units of fb. A smaller excluded cross-section implies a more stringent limit.*

Figure 6.37: *Exclusion contours and excluded cross-section limits at 95% confidence level in the parameter space of one-step simplified models with gluino pair production and decays according to $\tilde{g} \rightarrow q\bar{q}'\tilde{\chi}_1^\pm \rightarrow q\bar{q}'W^\pm\tilde{\chi}_1^0$. The free parameters of the model are the gluino mass and the x parameter, whereas the neutralino mass is fixed to 60 GeV. The limits are shown for the soft-lepton (top left), hard-lepton (top right), and multi-lepton signal regions (bottom left), as well as for the combination of all signal regions (bottom right). The band around the median expected limit shows the $\pm 1\sigma$ variations, including all uncertainties except theoretical uncertainties on the signal. The dotted lines around the observed limit correspond to the $\pm 1\sigma$ variations of the theoretical uncertainties of the signal samples. The grey numbers indicate the excluded cross-section in units of fb. A smaller excluded cross-section implies a more stringent limit.*

Figure 6.38: Exclusion contours and excluded cross-section limits at 95% confidence level in the parameter space of one-step simplified models with squark pair production and decays according to $\tilde{q}_L \rightarrow q' \tilde{\chi}_1^\pm \rightarrow q' W^\pm \tilde{\chi}_1^0$. The x -parameter is either set to 1/2 (left) or variable at fixed neutralino mass (right). The limits are shown for the combination of all signal regions. The band around the median expected limit shows the $\pm 1\sigma$ variations, including all uncertainties except theoretical uncertainties on the signal. The dotted lines around the observed limit correspond to the $\pm 1\sigma$ variations of the theoretical uncertainties of the signal samples. The grey numbers indicate the excluded cross-section in units of fb.

Figure 6.39: Exclusion contours and excluded cross-section limits at 95% confidence level in the parameter space of two-step simplified models with gluino pair production and decays according to $\tilde{g} \rightarrow q\bar{q}' \tilde{\chi}_1^\pm \rightarrow q\bar{q}' \ell^\pm \tilde{\nu}_L \rightarrow q\bar{q}' \ell^\pm \nu \tilde{\chi}_1^0$ or $\tilde{g} \rightarrow q\bar{q}' \tilde{\chi}_1^\pm \rightarrow q\bar{q}' \nu \tilde{\ell}_L^\pm \rightarrow q\bar{q}' \nu \ell^\pm \tilde{\chi}_1^0$ (left) or $\tilde{g} \rightarrow q\bar{q}' \tilde{\chi}_1^\pm \rightarrow q\bar{q}' W^{(*)\pm} \tilde{\chi}_2^0 \rightarrow W^{(*)\pm} Z^{(*)\pm} \tilde{\chi}_1^0$ (right). The limits are shown for the combination of all signal regions. The band around the median expected limit shows the $\pm 1\sigma$ variations, including all uncertainties except theoretical uncertainties on the signal. The dotted lines around the observed limit correspond to the $\pm 1\sigma$ variations of the theoretical uncertainties of the signal samples. The grey numbers indicate the excluded cross-section in units of fb.

Figure 6.40: Exclusion contours and excluded cross-section limits at 95% confidence level in the parameter space of two-step simplified models with squark pair production. Top left: both squarks decay via $\tilde{q}_L \rightarrow q' \tilde{\chi}_1^{\pm} \rightarrow q' \ell^{\pm} \tilde{\nu}_L \rightarrow q' \ell^{\pm} \nu \tilde{\chi}_1^0$ or $\tilde{q}_L \rightarrow q' \tilde{\chi}_1^{\pm} \rightarrow q' \tilde{\ell}^{\pm} \nu \rightarrow q' \ell^{\pm} \nu \tilde{\chi}_1^0$. Top right: one squark decays via $\tilde{q}_L \rightarrow q' \tilde{\chi}_1^{\pm} \rightarrow q' \ell^{\pm} \tilde{\nu}_L \rightarrow q' \ell^{\pm} \nu \tilde{\chi}_1^0$ or $\tilde{q}_L \rightarrow q' \tilde{\chi}_1^{\pm} \rightarrow q' \tilde{\ell}^{\pm} \nu \rightarrow q' \ell^{\pm} \nu \tilde{\chi}_1^0$ and the other squark decays via $\tilde{q}_L \rightarrow q \tilde{\chi}_2^0 \rightarrow q \ell^{\pm} \tilde{\ell}_L^{\mp} \rightarrow q \ell^{\pm} \ell^{\mp} \tilde{\chi}_1^0$ or $\tilde{q}_L \rightarrow q \tilde{\chi}_2^0 \rightarrow q \nu \tilde{\nu}_L \rightarrow q \nu \nu \tilde{\chi}_1^0$. Bottom row: both squarks decay via $\tilde{q}_L \rightarrow q' \tilde{\chi}_1^{\pm} \rightarrow W^{(*)\pm} \tilde{\chi}_2^0 \rightarrow W^{(*)\pm} Z^{(*)} \tilde{\chi}_1^0$. The limits are shown for the combination of all signal regions. The band around the median expected limit shows the $\pm 1\sigma$ variations, including all uncertainties except theoretical uncertainties on the signal. The dotted lines around the observed limit correspond to the $\pm 1\sigma$ variations of the theoretical uncertainties of the signal samples. The grey numbers indicate the excluded cross-section in units of fb. A smaller excluded cross-section implies a more stringent limit.

6.11 Conclusion

The first search at the LHC for compressed supersymmetric scenarios in final states with missing transverse energy, jets, and one isolated low- p_T lepton has been presented in this chapter using the full 2011 ATLAS dataset of 4.7 fb^{-1} in proton-proton collisions at 7 TeV center-of-mass energy. The events in the analysis are collected with missing transverse energy triggers. The event selection is based on the presence of initial state radiation jets to increase the signal acceptance and the discrimination power against Standard Model backgrounds. The dominant backgrounds, from W/Z +jets and $t\bar{t}$ processes, are measured in enriched control regions. The smaller multijet background is determined with an independent data-driven technique and the remaining minor backgrounds from di-boson, single-top, and $t\bar{t}$ plus vector boson processes are taken from simulation. The W/Z +jets simulated samples are reweighted in bins of generated vector boson p_T using a fit to the reconstructed p_T^Z spectrum to better describe the data. The background estimates and simulated samples are input to a global likelihood fit which uses the shape of the jet multiplicity distributions in the $t\bar{t}$ and W/Z +jets control regions. The estimated event yields in the control regions are extrapolated to the signal regions with transfer factors taken from simulation. The combined fit is overconstrained such that some systematic uncertainties are reduced in the fit using data. The validity of the background estimates is tested in validation regions that lie between the control and the signal regions. No significant excess over the Standard Model expectation is observed in any of the signal regions. Model independent upper limits on the visible signal cross-section are calculated. Exclusion limits for one-step and two-step simplified models are derived from an additional fit in the signal regions. The sensitivity in the compressed mass region of these models is improved by up to a factor of ten with respect to two complementary analyses with one or two high- p_T leptons in the final state. The results of all three analyses are combined statistically to provide the best exclusion limits over the full parameter range of the considered supersymmetric scenarios. The results of this analysis have been published in [166].

Chapter 7

Conclusion

Three years after the highly-anticipated start-up of the LHC at 7 TeV center-of-mass energy, the ATLAS experiment has performed a diverse program of physics measurements and searches at the energy frontier. Part of this program are the two physics analyses presented in this thesis, the measurement of the inclusive cross-section for electron production from heavy flavour decays using ATLAS data collected in 2010, and the search for compressed supersymmetric scenarios in events with missing transverse energy, jets, and one isolated low transverse momentum lepton in the final state based on data taken in 2011.

The measurement of the inclusive cross-section for electron production from heavy-flavour decays is motivated by the possibility to validate and constrain theoretical predictions on heavy flavour production in the new energy regime of the LHC. Such predictions are important for future measurements and searches where heavy-flavour processes constitute a major background. In addition, the analysis allows to evaluate the performance of the ATLAS electron reconstruction in the low transverse momentum range, where the spectrum is dominated by non-isolated electrons from heavy-flavour decays. A good understanding of the reconstruction performance in this regime forms the basis for new physics searches where low- p_T electrons are expected in the final state.

The cross-section measurement presented in this thesis was performed in the electron transverse momentum range $7 \text{ GeV} < p_T < 26 \text{ GeV}$ and within $|\eta| < 2.0$ in pseudorapidity, excluding the transition region $1.37 < |\eta| < 1.52$ between the barrel and endcap electromagnetic calorimeters. The measurement was not extended to higher transverse momenta since the contribution of isolated electrons from $W/Z/\gamma^*$ production processes, which have different reconstruction and identification efficiencies, becomes dominant in that region. The analysis is based on a data sample of 1.3 pb^{-1} , taken at 7 TeV center-of-mass energy in proton-proton collisions. The data were collected with electromagnetic calorimeter triggers during the early days of LHC operation in 2010, where the low instantaneous luminosity allowed to select low- p_T electrons without the application of trigger prescale factors or further high-level trigger requirements that would reduce the signal efficiency for non-isolated electrons or bias the selection. A three-dimensional extended binned maximum likelihood technique was used to extract the number of signal electrons from the data sample and to discriminate against large backgrounds from hadrons faking electrons and electrons from photon conversions. While the trigger efficiencies were determined from data, the signal selection efficiencies were derived from high statistics simulated samples. The uncertainties associated to the signal efficiencies were determined from independent data

measurements and simulated samples assuming additional material with respect to the nominal detector geometry. The measurement was unfolded to true electron transverse momentum using a bin-by-bin unfolding technique. Within the uncertainties, the measured cross-section was found to be compatible with the state-of-the-art theoretical prediction of the FONLL framework at NLO+NLL accuracy. In addition, the measurement was found to be consistent with the POWHEG+PYTHIA prediction, and with a complementary ATLAS measurement with muons in the final state. The total systematic uncertainty of the measurement ranges between 15-18 %, with the largest contribution coming from the determination of the efficiency and migration correction. The measurement was published in [163] together with the complementary inclusive muon analysis.

The search for compressed supersymmetric scenarios presented in the second part of this thesis is motivated by the need to explore experimentally less accessible regions of supersymmetric parameter space given the existing stringent bounds on SUSY. In compressed supersymmetric scenarios the predicted superparticles are close in mass thus leading to experimental signatures with soft physics objects in the final state. The identification of such low transverse momentum objects is exacerbated by the presence of large backgrounds in this regime.

The search presented in this work is based on events with missing transverse energy, jets, and one low transverse momentum lepton in the final state. The analysis thus builds, in part, on the previously described measurement using low transverse momentum electrons. The data sample used in the analysis amounts to 4.7 fb^{-1} and was collected in 2011 at 7 TeV center-of-mass energy in proton-proton collisions using missing transverse energy triggers. The event selection is based on the presence of initial state radiation jets to increase the signal acceptance and the discrimination power against the Standard Model backgrounds. The dominant backgrounds, from W/Z +jets and $t\bar{t}$ processes, were isolated and measured in control regions. The smaller multijet background was determined with an independent data-driven technique and the remaining minor backgrounds from di-boson, single-top, and $t\bar{t}$ plus vector boson processes were taken from simulation. The W/Z +jets simulated samples were reweighted in bins of generated vector boson p_T using a fit to the reconstructed p_T^Z spectrum to better describe the data. The background estimates and simulated samples were used as input to a global likelihood fit based on the shape of the jet multiplicity distributions in the $t\bar{t}$ and W/Z +jets control regions. The estimated event yields in the control regions were extrapolated to the signal regions with transfer factors taken from simulation. The validity of the background estimates was tested in validation regions that lie between the control and the signal regions. No significant excess over the Standard Model expectation was observed in any of the signal regions. Model independent upper limits on the visible signal cross-section were calculated. Exclusion limits for one-step and two-step simplified models were derived from an additional fit in the signal regions. The sensitivity in the compressed mass region of these models was found to be improved by up to a factor of ten with respect to two complementary analyses with one or two high- p_T leptons in the final state. The results of all three analyses were combined statistically to provide the best exclusion limits over the full parameter range of the considered supersymmetric scenarios. The results of this analysis were published in [166] and constitute the first search at the LHC for compressed SUSY with soft leptons in the final state.

Appendix A

Supplementary Information for Chapter 5

A.1 Supplementary Information on the Signal Extraction

E_T bin [GeV]	7-8	8-10	10-12	12-14	14-16	16-18	18-20	20-22	22-26
Corrected	2934±127	2982±108	10337±313	4912±180	6000±163	3174±116	7490±236	4757±154	5203±186
Scaled up	3060±132	3126±113	10572±341	5077±190	6181 174	3284±122	7328±252	4925±163	5339±191
Scaled down	2933±127	2972±108	10019±294	4805±171	5714 157	3143±111	7514±220	4658±146	5024±180

Table A.1: *Summary of $N^{Q \rightarrow e}$ signal extraction results obtained with the energy scale corrected and the associated uncertainty shifted up and down.*

E_T -range [GeV]	$ \eta < 0.625$	$0.625 < \eta < 1.07$	$1.07 < \eta < 1.37$	$1.52 < \eta < 2.0$	Combined η
$7 < E_T < 8$	1.0	1.5	3.7	2.4	0.9
$8 < E_T < 10$	1.0	1.6	2.1	2.6	0.8
$10 < E_T < 12$	1.6	2.2	3.2	5.1	1.4
$12 < E_T < 14$	1.9	3.6	5.0	5.5	1.8
$14 < E_T < 16$	2.7	3.9	5.7	9.8	2.5
$16 < E_T < 18$	1.6	2.5	4.2	4.4	1.4
$18 < E_T < 20$	1.9	2.7	6.1	8.8	2.0
$20 < E_T < 22$	3.0	3.6	4.7	14.1	2.5
$22 < E_T < 26$	2.4	3.4	5.2	13.1	2.5

Table A.2: *Systematic uncertainties (in percent) on the extracted $N^{Q \rightarrow e}$ yield due to the finite statistics of the MC samples used to create the pdfs. The uncertainties are shown in bins of E_T and $|\eta|$ and for the combination over all $|\eta|$ bins taking into account correlations between bins. Since different MC samples are used in different E_T regions this uncertainty has a genuine E_T dependence which is not further averaged.*

A.2 Supplementary Information on the Efficiency Determination and Unfolding Procedure

Figure A.1: *Left: Comparison of the efficiencies and migration corrections as well as their ratio for PYTHIA and POWHEG simulated samples. Right: Ratio of the combined efficiency and migration correction factors of the POWHEG simulated samples with respect to the PYTHIA simulated samples. Taken from [134].*

E_T -range [GeV]	7-8	8-10	10-12	12-14	14-16	16-18	18-20	20-22	22-26
Input: PYTHIA									
Truth	695.7 $\pm$ 1.8	649.1 $\pm$ 1.7	258.9 $\pm$ 1.1	115.9 $\pm$ 0.7	57.4 $\pm$ 0.5	30.7 $\pm$ 0.4	17.6 $\pm$ 0.3	10.6 $\pm$ 0.2	10.1 $\pm$ 0.2
Bin-by-bin	695.7 $\pm$ 6.1	649.1 $\pm$ 5.1	258.9 $\pm$ 3.2	115.9 $\pm$ 2.1	57.4 $\pm$ 1.4	30.7 $\pm$ 1.1	17.6 $\pm$ 0.8	10.6 $\pm$ 0.6	10.1 $\pm$ 0.6
w. toy MC	695.7 $\pm$ 4.3	649.1 $\pm$ 3.7	258.9 $\pm$ 1.9	115.9 $\pm$ 1.5	57.4 $\pm$ 1.1	30.7 $\pm$ 0.8	17.6 $\pm$ 0.6	10.6 $\pm$ 0.5	10.1 $\pm$ 0.4
Bayesian (n=4)	711.5 $\pm$ 5.3	650.8 $\pm$ 4.3	239.0 $\pm$ 2.6	113.8 $\pm$ 1.9	59.5 $\pm$ 1.3	31.6 $\pm$ 1.0	18.5 $\pm$ 0.6	11.0 $\pm$ 0.5	14.7 $\pm$ 0.6
Bayesian (n=1)	704.1 $\pm$ 3.1	648.4 $\pm$ 2.1	248.9 $\pm$ 1.2	113.7 $\pm$ 1.0	58.2 $\pm$ 0.8	31.7 $\pm$ 0.5	18.7 $\pm$ 0.4	11.7 $\pm$ 0.3	13.1 $\pm$ 0.4
SVD (k=4)	714.4 $\pm$ 3.6	642.9 $\pm$ 3.1	247.3 $\pm$ 1.5	112.3 $\pm$ 0.9	57.7 $\pm$ 0.6	32.3 $\pm$ 0.4	19.4 $\pm$ 0.3	12.2 $\pm$ 0.2	12.1 $\pm$ 0.3
SVD (k=2)	695.7 $\pm$ 3.2	648.4 $\pm$ 2.7	258.5 $\pm$ 1.2	115.9 $\pm$ 0.6	57.5 $\pm$ 0.4	30.9 $\pm$ 0.2	17.7 $\pm$ 0.2	10.7 $\pm$ 0.1	10.3 $\pm$ 0.1
Input: POWHEG+PYTHIA									
Truth	309.3 $\pm$ 1.4	296.8 $\pm$ 1.4	121.2 $\pm$ 0.9	57.4 $\pm$ 0.6	28.1 $\pm$ 0.4	16.3 $\pm$ 0.3	9.0 $\pm$ 0.2	6.0 $\pm$ 0.2	6.0 $\pm$ 0.2
Bin-by-bin	307.1 $\pm$ 3.1	294.7 $\pm$ 2.7	123.3 $\pm$ 1.8	57.2 $\pm$ 1.2	28.6 $\pm$ 0.8	16.9 $\pm$ 0.7	8.9 $\pm$ 0.5	5.9 $\pm$ 0.4	5.6 $\pm$ 0.4
w. toy MC	307.1 $\pm$ 2.4	294.7 $\pm$ 2.1	123.3 $\pm$ 1.1	57.2 $\pm$ 1.0	28.6 $\pm$ 0.7	16.9 $\pm$ 0.6	8.9 $\pm$ 0.4	5.9 $\pm$ 0.3	5.6 $\pm$ 0.3
Bayesian (n=4)	311.6 $\pm$ 3.0	295.5 $\pm$ 2.5	115.0 $\pm$ 1.6	56.4 $\pm$ 1.2	29.7 $\pm$ 0.8	17.5 $\pm$ 0.7	9.5 $\pm$ 0.4	6.0 $\pm$ 0.3	8.1 $\pm$ 0.4
Bayesian (n=1)	313.3 $\pm$ 1.7	294.5 $\pm$ 1.3	117.7 $\pm$ 0.7	55.6 $\pm$ 0.6	29.2 $\pm$ 0.5	16.7 $\pm$ 0.3	9.7 $\pm$ 0.2	6.3 $\pm$ 0.2	7.2 $\pm$ 0.3
SVD (k=4)	312.0 $\pm$ 2.0	293.2 $\pm$ 1.9	118.1 $\pm$ 0.9	55.7 $\pm$ 0.6	29.5 $\pm$ 0.4	17.0 $\pm$ 0.2	10.4 $\pm$ 0.2	6.6 $\pm$ 0.2	6.6 $\pm$ 0.2
SVD (k=2)	312.2 $\pm$ 1.9	296.6 $\pm$ 1.6	120.8 $\pm$ 0.7	55.3 $\pm$ 0.4	27.9 $\pm$ 0.2	15.3 $\pm$ 0.2	8.9 $\pm$ 0.1	5.4 $\pm$ 0.1	5.2 $\pm$ 0.1
Input: POWHEG+JIMMY+HERWIG									
Truth	186.3 $\pm$ 0.9	173.3 $\pm$ 0.9	70.7 $\pm$ 0.6	32.3 $\pm$ 0.4	16.7 $\pm$ 0.3	9.1 $\pm$ 0.2	5.1 $\pm$ 0.2	3.0 $\pm$ 0.1	3.2 $\pm$ 0.1
Bin-by-bin	181.8 $\pm$ 1.9	172.9 $\pm$ 1.6	71.2 $\pm$ 1.1	32.8 $\pm$ 0.7	16.7 $\pm$ 0.5	9.3 $\pm$ 0.4	5.2 $\pm$ 0.3	3.2 $\pm$ 0.2	3.2 $\pm$ 0.2
w. toy MC	181.8 $\pm$ 1.5	172.9 $\pm$ 1.3	71.2 $\pm$ 0.7	32.8 $\pm$ 0.6	16.7 $\pm$ 0.4	9.3 $\pm$ 0.3	5.2 $\pm$ 0.2	3.2 $\pm$ 0.2	3.2 $\pm$ 0.2
Bayesian (n=4)	184.9 $\pm$ 1.9	173.4 $\pm$ 1.6	66.1 $\pm$ 1.0	32.4 $\pm$ 0.7	17.3 $\pm$ 0.5	9.6 $\pm$ 0.4	5.5 $\pm$ 0.2	3.3 $\pm$ 0.2	4.6 $\pm$ 0.2
Bayesian (n=1)	185.0 $\pm$ 1.1	172.7 $\pm$ 0.8	68.1 $\pm$ 0.4	32.0 $\pm$ 0.4	16.8 $\pm$ 0.3	9.4 $\pm$ 0.2	5.5 $\pm$ 0.1	3.5 $\pm$ 0.1	4.1 $\pm$ 0.2
SVD (k=4)	185.3 $\pm$ 1.3	171.7 $\pm$ 1.2	68.2 $\pm$ 0.5	32.0 $\pm$ 0.4	16.9 $\pm$ 0.2	9.7 $\pm$ 0.1	5.9 $\pm$ 0.1	3.7 $\pm$ 0.1	3.7 $\pm$ 0.1
SVD (k=2)	183.9 $\pm$ 1.2	173.8 $\pm$ 1.0	70.4 $\pm$ 0.4	32.0 $\pm$ 0.2	16.1 $\pm$ 0.2	8.8 $\pm$ 0.1	5.1 $\pm$ 0.1	3.1 $\pm$ 0.0	3.0 $\pm$ 0.0
Input: PYTHIA distorted material									
Truth	693.2 $\pm$ 3.5	655.7 $\pm$ 3.4	259.3 $\pm$ 2.2	115.9 $\pm$ 1.4	58.8 $\pm$ 1.0	30.3 $\pm$ 0.7	17.7 $\pm$ 0.6	11.1 $\pm$ 0.4	10.5 $\pm$ 0.4
Bin-by-bin	598.6 $\pm$ 6.4	569.0 $\pm$ 5.7	229.0 $\pm$ 3.6	103.1 $\pm$ 2.4	54.4 $\pm$ 1.7	26.7 $\pm$ 1.2	15.8 $\pm$ 0.9	9.5 $\pm$ 0.7	9.3 $\pm$ 0.7
w. toy MC	598.6 $\pm$ 5.3	569.0 $\pm$ 4.7	229.0 $\pm$ 2.4	103.1 $\pm$ 2.0	54.4 $\pm$ 1.5	26.7 $\pm$ 1.1	15.8 $\pm$ 0.8	9.5 $\pm$ 0.6	9.3 $\pm$ 0.5
Bayesian (n=4)	609.7 $\pm$ 6.6	571.4 $\pm$ 5.6	211.8 $\pm$ 3.4	101.1 $\pm$ 2.5	57.1 $\pm$ 1.8	27.6 $\pm$ 1.3	16.4 $\pm$ 0.8	9.9 $\pm$ 0.7	13.5 $\pm$ 0.8
Bayesian (n=1)	608.8 $\pm$ 3.8	567.0 $\pm$ 2.8	219.9 $\pm$ 1.5	101.6 $\pm$ 1.3	53.7 $\pm$ 1.0	28.4 $\pm$ 0.6	16.7 $\pm$ 0.5	10.5 $\pm$ 0.4	12.0 $\pm$ 0.5
SVD (k=4)	612.6 $\pm$ 4.5	563.3 $\pm$ 4.0	219.8 $\pm$ 1.9	101.2 $\pm$ 1.2	52.5 $\pm$ 0.8	29.3 $\pm$ 0.5	17.5 $\pm$ 0.4	11.0 $\pm$ 0.3	10.9 $\pm$ 0.3
SVD (k=2)	604.9 $\pm$ 4.0	567.2 $\pm$ 3.5	227.6 $\pm$ 1.5	102.7 $\pm$ 0.8	51.3 $\pm$ 0.5	27.7 $\pm$ 0.3	15.9 $\pm$ 0.2	9.6 $\pm$ 0.1	9.3 $\pm$ 0.1
Input: Extracted electron data									
Bin-by-bin	351.2 $\pm$ 15.5	334.3 $\pm$ 12.3	134.0 $\pm$ 4.5	60.5 $\pm$ 2.6	30.6 $\pm$ 1.2	16.1 $\pm$ 0.8	9.2 $\pm$ 0.5	5.5 $\pm$ 0.4	5.2 $\pm$ 0.4
w. toy MC	351.2 $\pm$ 15.5	334.3 $\pm$ 11.9	134.0 $\pm$ 3.5	60.5 $\pm$ 2.5	30.6 $\pm$ 1.1	16.1 $\pm$ 0.8	9.2 $\pm$ 0.4	5.5 $\pm$ 0.3	5.2 $\pm$ 0.3
Bayesian (n=4)	357.7 $\pm$ 19.1	335.9 $\pm$ 14.0	124.0 $\pm$ 4.9	59.4 $\pm$ 3.2	31.9 $\pm$ 1.3	16.6 $\pm$ 1.0	9.6 $\pm$ 0.4	5.7 $\pm$ 0.3	7.5 $\pm$ 0.4
Bayesian (n=1)	357.3 $\pm$ 11.2	332.9 $\pm$ 7.0	128.8 $\pm$ 2.3	59.3 $\pm$ 1.6	30.8 $\pm$ 0.8	16.7 $\pm$ 0.5	9.7 $\pm$ 0.3	6.1 $\pm$ 0.2	6.8 $\pm$ 0.2
SVD (k=4)	357.2 $\pm$ 11.1	327.3 $\pm$ 6.7	129.3 $\pm$ 2.5	59.4 $\pm$ 1.4	30.5 $\pm$ 0.7	16.8 $\pm$ 0.4	10.0 $\pm$ 0.2	6.4 $\pm$ 0.2	6.5 $\pm$ 0.2
SVD (k=2)	354.0 $\pm$ 7.3	334.1 $\pm$ 6.3	135.0 $\pm$ 2.4	61.4 $\pm$ 1.1	30.9 $\pm$ 0.5	16.8 $\pm$ 0.3	9.7 $\pm$ 0.2	5.9 $\pm$ 0.1	5.7 $\pm$ 0.1

Table A.3: Comparison of the performance of the different unfolding methods. For SVD unfolding, k is the regularization parameter, while for Bayesian unfolding, n is the number of iterations. In all cases PYTHIA simulated samples of electrons from heavy flavour decays are used to build the response matrix. The first block is a closure test where the same Monte-Carlo sample is used to build the response matrix and the input. The next three blocks show the bias of the methods on independent MC samples (POWHEG+PYTHIA, POWHEG+JIMMY+HERWIG, and PYTHIA with extra material). The last block shows the results on the data. Only statistical errors are shown. Taken from [134].

Appendix B

Supplementary Information for Chapter 6

B.1 Results in the Hard-lepton and Multi-lepton Validation Regions

Electron channel	High- m_T VR	W VR	$t\bar{t}$ VR
Observed	7698	1068	499
Fitted bkg	7618.1 ± 553.9	1017.8 ± 42.9	501.6 ± 16.7
Fitted top	3014.6 ± 113.3	213.5 ± 16.9	410.6 ± 17.5
Fitted W/Z +jets	3936.4 ± 162.7	778.1 ± 39.5	58.2 ± 9.7
Fitted other bkg	231.2 ± 46.3	26.2 ± 5.7	32.8 ± 6.6
Fitted multijet	435.9 ± 525.6	0.0 ± 0.0	0.0 ± 0.0
MC exp. SM	8054.7	1160.3	561.7
MC exp. top	3382.1	281.0	447.7
MC exp. W/Z +jets	3987.6	842.3	72.7
MC exp. other bkg	249.1	37.0	41.3
Data-driven multijet	435.9	0.0	0.0

Table B.1: *Observed and expected event yields in the hard-electron validation regions. The fitted number of events contain statistical and systematic uncertainties. In addition the inputs to the fit from simulation as well as the data-driven multijet estimate are shown.*

Muon channel	High-m_T VR	W VR	$t\bar{t}$ VR
Observed	7088	1020	428
Fitted bkg	7298.0 ± 152.8	1038.0 ± 37.7	474.9 ± 11.6
Fitted top	2884.0 ± 111.9	192.1 ± 14.5	371.1 ± 14.8
Fitted W/Z +jets	4148.8 ± 143.5	816.9 ± 34.0	74.6 ± 9.5
Fitted other bkg	229.6 ± 45.8	24.1 ± 5.1	26.1 ± 5.3
Fitted multijet	35.7 ± 85.0	5.0 ± 11.5	3.2 ± 5.2
MC exp. SM	7503.2	1120.9	493.5
MC exp. top	3213.3	243.2	386.8
MC exp. W/Z +jets	3998.6	840.7	74.7
MC exp. other bkg	255.7	32.0	28.8
Data-driven multijet	35.7	5.0	3.2

Table B.2: Observed and expected event yields in the hard-muon validation regions. The fitted number of events contain statistical and systematic uncertainties. In addition the inputs to the fit from simulation as well as the data-driven multijet estimate are shown.

Multi-lepton	2-jet $t\bar{t}$ VR			2-jet Z VR		
	el-el	mu-mu	el-mu	el-el	mu-mu	el-mu
Observed	64	83	117	150	181	34
Fitted bkg	53.4 ± 3.7	74.7 ± 4.8	110.9 ± 7.8	134.4 ± 7.3	173.3 ± 8.7	30.4 ± 3.5
Fitted top	39.8 ± 2.9	54.5 ± 4.1	100.4 ± 7.2	7.7 ± 1.3	10.5 ± 1.3	22.0 ± 2.8
Fitted W/Z +jets	11.5 ± 0.8	17.8 ± 1.0	5.3 ± 0.4	125.6 ± 6.2	161.6 ± 8.3	6.0 ± 0.5
Fitted other bkg	2.1 ± 0.4	2.5 ± 0.5	5.3 ± 1.1	1.1 ± 0.3	1.2 ± 0.3	1.2 ± 0.3
Fitted multijet	0.0 ± 0.0	0.0 ± 0.0	0.0 ± 0.0	0.0 ± 0.0	0.0 ± 0.0	1.2 ± 1.5
MC exp. SM	61.3	85.5	134.5	137.3	176.5	42.6
MC exp. top	47.1	65.5	122.8	12.3	16.1	33.5
MC exp. W/Z +jets	12.0	17.2	5.6	123.8	158.9	6.6
MC exp. other bkg	2.2	2.8	6.1	1.3	1.5	1.3
Data-driven multijet	0.0	0.0	0.0	0.0	0.0	1.2

Table B.3: Observed and expected event yields in the 2-jet multi-lepton validation regions. The fitted number of events contain statistical and systematic uncertainties. In addition the inputs to the fit from simulation as well as the data-driven multijet estimate are shown.

Multi-lepton	4-jet $t\bar{t}$ VR			4-jet Z VR		
	el-el	mu-mu	el-mu	el-el	mu-mu	el-mu
Observed	8	12	22	40	34	23
Fitted bkg	9.4 ± 1.2	11.0 ± 1.0	20.0 ± 1.5	33.1 ± 1.9	36.8 ± 2.2	24.2 ± 2.0
Fitted top	8.6 ± 0.9	10.2 ± 1.0	19.7 ± 1.5	8.6 ± 0.7	12.5 ± 0.9	22.3 ± 1.8
Fitted W/Z +jets	0.3 ± 0.1	0.4 ± 0.1	0.0 ± 0.0	24.0 ± 1.4	24.1 ± 1.6	1.5 ± 0.2
Fitted other bkg	0.3 ± 0.1	0.3 ± 0.1	0.3 ± 0.1	0.2 ± 0.1	0.2 ± 0.0	0.0 ± 0.0
Fitted multijet	0.2 ± 0.6	0.0 ± 0.0	0.0 ± 0.0	0.2 ± 0.6	0.0 ± 0.0	0.4 ± 0.8
MC exp. SM	10.3	12.7	23.0	34.5	39.7	31.8
MC exp. top	9.6	12.0	22.7	11.4	15.9	29.6
MC exp. W/Z +jets	0.3	0.3	0.0	22.6	23.5	1.6
MC exp. other bkg	0.2	0.4	0.3	0.4	0.3	0.3
Data-driven multijet	0.2	0.0	0.0	0.2	0.0	0.4

Table B.4: Observed and expected event yields in the 4-jet multi-lepton validation regions. The fitted number of events contain statistical and systematic uncertainties. In addition the inputs to the fit from simulation as well as the data-driven multijet estimate are shown.

B.2 Results in the Hard-lepton and Multi-lepton Signal Regions

Single-lepton	Electron			Muon		
	3-jet	4-jet	soft lepton	3-jet	4-jet	soft lepton
Observed	2	4	11	1	2	14
Fitted bkg	2.3 ± 0.9	3.5 ± 0.9	14.0 ± 3.3	2.6 ± 0.8	1.5 ± 0.3	19 ± 5
Fitted top	0.4 ± 0.2	2.3 ± 0.6	3.8 ± 0.6	0.5 ± 0.2	1.3 ± 0.3	3.8 ± 0.8
Fitted W/Z +jets	1.5 ± 0.6	0.9 ± 0.2	5.8 ± 1.0	2.0 ± 0.6	0.2 ± 0.1	11.4 ± 2.3
Fitted other bkg	0.0 ± 0.0	$0.0^{+0.3}_{-0.0}$	0.6 ± 0.1	0.1 ± 0.1	0.0 ± 0.0	0.2 ± 0.1
Fitted multijet	0.3 ± 0.4	0.3 ± 0.4	3.8 ± 2.5	0.0 ± 0.0	0.0 ± 0.0	3.6 ± 2.5
MC exp. SM	2.7	5.3	14.2	2.8	2.4	18.0
MC exp. top	0.9	3.1	4.3	0.6	2.0	3.8
MC exp. W/Z +jets	1.5	1.3	5.5	2.0	0.3	10.5
MC exp. other bkg	0.0	0.5	0.5	0.2	0.1	0.1
Data-driven multijet	0.3	0.3	3.8	0.0	0.0	3.6

Table B.5: Observed and expected event yields in the soft-lepton and hard-lepton signal regions. The fitted number of events contain statistical and systematic uncertainties. In addition the inputs to the fit from simulation as well as the data-driven multijet estimate are shown.

Multi-lepton	2-jets			4-jets		
	ee	$\mu\mu$	$e\mu$	ee	$\mu\mu$	$e\mu$
Observed	0	0	1	8	12	18
Fitted bkg	0.3 ± 0.2	0.4 ± 0.2	0.7 ± 0.2	9.1 ± 1.5	11.7 ± 1.7	21 ± 3
Fitted top	0.1 ± 0.1	0.2 ± 0.1	0.6 ± 0.2	9.1 ± 1.4	11.1 ± 1.7	20 ± 3
Fitted W/Z +jets	0.1 ± 0.1	0.1 ± 0.0	0.0 ± 0.0	0.0 ± 0.0	0.2 ± 0.1	0.4 ± 0.1
Fitted other bkg	0.1 ± 0.1	0.1 ± 0.0	0.1 ± 0.0	0.0 ± 0.0	0.4 ± 0.1	0.6 ± 0.1
Fitted multijet	0.0 ± 0.0	0.0 ± 0.0	0.0 ± 0.0	0.0 ± 0.2	0.0 ± 0.0	0.0 ± 0.0
MC exp. SM	0.3	0.5	0.9	11.4	14.7	27.1
MC exp. top	0.2	0.3	0.7	11.1	13.9	26.0
MC exp. W/Z +jets	0.1	0.1	0.1	0.1	0.3	0.4
MC exp. other bkg	0.1	0.1	0.1	0.2	0.5	0.7
Data-driven multijet	0.0	0.0	0.0	0.0	0.0	0.0

Table B.6: Observed and expected event yields in the multi-lepton signal regions. The fitted number of events contain statistical and systematic uncertainties. In addition the inputs to the fit from simulation as well as the data-driven multijet estimate are shown.

B.3 Visible Cross-Section Upper Limits in the Hard-lepton and Multi-lepton Analyses

Signal channel	$\langle\epsilon\sigma\rangle_{\text{obs}}^{95}[\text{fb}]$	S_{obs}^{95}	S_{exp}^{95}	CL_B
soft electron	1.82	8.6	$10.4^{+4.2}_{-3.1}$	0.28
soft muon	1.92	9.0	$12.5^{+5.4}_{-3.8}$	0.21
hard electron, 3-jet	0.94	4.4	$4.3^{+2.0}_{-0.8}$	0.54
hard muon, 3-jet	0.75	3.6	$4.2^{+2.0}_{-0.7}$	0.27
hard electron, 4-jet	1.22	5.8	$5.3^{+2.6}_{-1.3}$	0.63
hard muon, 4-jet	0.95	4.5	$3.8^{+1.3}_{-0.7}$	0.75
multi-lepton, ee , 2-jet	0.71	3.3	3.5 ± 0.1	0.48
multi-lepton, $\mu\mu$, 2-jet	0.76	3.6	3.5 ± 0.1	0.46
multi-lepton, $e\mu$, 2-jet	0.83	3.9	$3.6^{+0.6}_{-0.2}$	0.85
multi-lepton, ee , 4-jet	1.53	7.2	$7.7^{+3.2}_{-2.1}$	0.39
multi-lepton, $\mu\mu$, 4-jet	1.93	9.1	$8.8^{+3.3}_{-3.0}$	0.55
multi-lepton, $e\mu$, 4-jet	2.14	10.1	$11.5^{+4.8}_{-3.5}$	0.28

Table B.7: Observed 95% CL upper limits on the visible cross section ($\langle\epsilon\sigma\rangle_{\text{obs}}^{95}$) as well as observed (S_{obs}^{95}) and expected (S_{exp}^{95}) upper limits on the number of signal events in the soft-lepton, hard-lepton, and multi-lepton signal regions. In addition the CL_b value, i.e. the observed confidence level for the background-only hypothesis, is shown.

B.4 Acceptance $\times$ Selection Efficiency for Signal Models

Figure B.1: Acceptance $\times$ efficiency in the soft-lepton (a)-(b) and hard-lepton (c)-(f) signal regions as a function of gluino and neutralino mass for a one-step simplified model with gluino production and mass ratio parameter $x = 1/2$.

Figure B.2: Acceptance $\times$ efficiency in the multi-lepton signal regions as a function of gluino and neutralino mass for a one-step simplified model with gluino production and mass ratio parameter $x = 1/2$.

Figure B.3: *Acceptance $\times$ efficiency in the soft-lepton (a)-(b) and hard-lepton (c)-(f) signal regions as a function of squark and neutralino mass for a one-step simplified model with squark production and mass ratio parameter $x = 1/2$.*

Figure B.4: Acceptance $\times$ efficiency in the multi-lepton signal regions as a function of squark and neutralino mass for a one-step simplified model with squark production and mass ratio parameter $x = 1/2$.

Figure B.5: *Acceptance $\times$ efficiency in the soft-lepton (a)-(b) and hard-lepton (c)-(f) signal regions as a function of gluino mass and mass ratio parameter for a one-step simplified model with gluino production and variable mass ratio parameter and fixed neutralino mass.*

Figure B.6: Acceptance $\times$ efficiency in the multi-lepton signal regions as a function of gluino mass and mass ratio parameter for a one-step simplified model with gluino production and variable mass ratio parameter and fixed neutralino mass.

Figure B.7: *Acceptance $\times$ efficiency in the soft-lepton (a)-(b) and hard-lepton (c)-(f) signal regions as a function of gluino mass and mass ratio parameter for a one-step simplified model with gluino production and variable mass ratio parameter and fixed neutralino mass.*

Figure B.8: Acceptance $\times$ efficiency in the multi-lepton signal regions as a function of gluino mass and mass ratio parameter for a one-step simplified model with gluino production and variable mass ratio parameter and fixed neutralino mass.

Figure B.9: Acceptance $\times$ efficiency in the soft-lepton (a)-(b) and hard-lepton (c)-(f) signal regions as a function of gluino and neutralino mass for a two-step simplified model with gluino production and decays as indicated in the figures.

Figure B.10: Acceptance $\times$ efficiency in the multi-lepton signal regions as a function of gluino and neutralino mass for a two-step simplified model with gluino production and decays as indicated in the figures.

Figure B.11: Acceptance $\times$ efficiency in the soft-lepton (a)-(b) and hard-lepton (c)-(f) signal regions as a function of gluino and neutralino mass for a two-step simplified model with gluino production and decays as indicated in the figures.

Figure B.12: Acceptance $\times$ efficiency in the multi-lepton signal regions as a function of gluino and neutralino mass for a two-step simplified model with gluino production and decays as indicated in the figures.

Figure B.13: Acceptance $\times$ efficiency in the soft-lepton (a)-(b) and hard-lepton (c)-(f) signal regions as a function of squark and neutralino mass for a two-step simplified model with squark production and decays as indicated in the figures.

Figure B.14: Acceptance $\times$ efficiency in the multi-lepton signal regions as a function of squark and neutralino mass for a two-step simplified model with squark production and decays as indicated in the figures.

Figure B.15: Acceptance $\times$ efficiency in the soft-lepton (a)-(b) and hard-lepton (c)-(f) signal regions as a function of squark and neutralino mass for a two-step simplified model with squark production and decays as indicated in the figures.

Figure B.16: Acceptance $\times$ efficiency in the multi-lepton signal regions as a function of squark and neutralino mass for a two-step simplified model with squark production and decays as indicated in the figures.

Figure B.17: Acceptance $\times$ efficiency in the soft-lepton (a)-(b) and hard-lepton (c)-(f) signal regions as a function of squark and neutralino mass for a two-step simplified model with squark production and decays as indicated in the figures.

Figure B.18: Acceptance $\times$ efficiency in the multi-lepton signal regions as a function of squark and neutralino mass for a two-step simplified model with squark production and decays as indicated in the figures.

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