SEARCH FOR CONTACT INTERACTIONS USING DIJET ANGULAR DISTRIBUTIONS WITH THE ATLAS DETECTOR AT THE CERN LARGE HADRON COLLIDER

by

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Abstract

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The LHC, with its center-of-mass energy of 7 TeV, offers the chance to investigate the fundamental constituents of matter at a higher energy scale than ever before. Using the data acquired by the ATLAS detector in the summer of 2010, two different measures of the angular distributions of dijet final states are studied and compared to Standard Model QCD expectations. Such a comparison is used to set new stringent limits on the existence of quark substructure.
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Chapter 1

Personal Contributions

The analysis summarized in this thesis was carried out by a collaboration of physicists participating on the ATLAS experiment. For the fulfillment of his doctoral degree, the author has made significant contributions to this analysis. For the sake of clarity, a list of the direct contributions to the analysis by the author is presented here:

1. The author carried out the full analysis of the $R_C$ observable presented in this thesis, which included:
   
   (a) Early truth Monte Carlo feasibility studies, concerned mostly with the effect of trigger biases and systematic uncertainties on the observable

   (b) Measurement of the $R_C$ distributions from both the ATLAS data and the fully simulated Monte Carlo samples, production of the associated plots, and quantification of the agreement between the Standard Model QCD predictions and the observed data

   (c) Generation of the NLO predictions for the $R_C$ observable, and estimation of the effects due to the associated variations in PDF error members and scale choices

   (d) Elaboration of the pseudo-experiment software framework which was used to
determine the effect of the convolution of the systematic uncertainties on the observable

(e) Implementation of the software used in setting Bayesian limits on the existence of contact interactions, along with the establishment of the associated technique used to interpolate between discrete contact interaction samples

2. The author formulated the requests for the official Monte Carlo contact interaction samples used in the limit setting procedure for both observables, along with the necessary associated validation tasks

3. The author implemented the software for Bayesian limit setting for the $\chi$ observable

Furthermore, for the sake of completeness, the author’s general contributions to the ATLAS experiment are listed here:

1. The author was in charge of maintaining the jet reconstruction software for a total of two years, a time period which overlapped with the collection of the first ATLAS data. The associated tasks included:

   (a) Implementing new features crucial to the preparation for first data, including a new calibration framework, routines for jet cleaning, and a correction for the vertex displacement

   (b) Providing support for software users, either directly, or through documentation and tutorials

   (c) Re-writing key parts of the software framework to improve performance and ease-of-use

2. The author participated in the effort for setting the jet energy scale uncertainty for first ATLAS data by studying the datasets with systematic variations, and implementing routines which are currently used to automate these studies
3. The author was in charge of the validation of Monte Carlo samples for the Jet + $E_T^{\text{miss}}$ group for a period of one year.

4. The author conducted a preliminary study of the feasibility of performing a search for leptoquarks using a data-driven template-based method in the dijet + $E_T^{\text{miss}}$ channel.
Chapter 2

Introduction

The field of particle physics has always been concerned with the study of the fundamental constituents of matter, and the forces that govern their interactions. This study has led to the establishment of the Standard Model (SM) of particle physics. The Standard Model is a quantum field theory consisting of twelve fermionic matter particles (along with their corresponding anti-particles) divided in two types (quarks and leptons), twelve bosonic force carriers, and a scalar boson (the Higgs) [1]. The fermions are further divided into 3 generations. The three fundamental forces, the electromagnetic, weak, and strong force, are mediated by the 8 gluons and 4 electroweak bosons ($Z, \gamma, W^+, W^-$).

The model obeys an $SU(3) \times SU(2) \times U(1)$ gauge symmetry, which dictates the coupling of the matter fields, through their gauge representation, to the corresponding gauge bosons. With the exception of the Higgs boson, all particles of the Standard Model have been experimentally observed. Of particular interest in the study which will be presented here is the strong force. It mediates the interactions between the quarks (which behave as triplets under $SU(3)$) and gluons (the $SU(3)$ gauge fields), described by the theory of Quantum ChromoDynamics (QCD). The quarks and gluons are colour-charged.
2.1 Quark Compositeness

Just as the atom was found to be divisible and composed of protons, neutrons and electrons, it has been postulated that the Standard Model quarks may be composite particles. While the particles which define the Standard Model are assumed to be indivisible, this assumption has only been verified up to a given length scale. The access to finer probes of matter (in the case of particle physics, higher energy regimes) provides an opportunity to explore the possibility of the presence of a finer substructure in the particles of the model. The LHC, which ran at a center-of-mass energy of 7 TeV in the 2010 run, offers a new window into a higher kinematic regime than ever studied at a particle collider beforehand, providing such an opportunity.

QCD dijet production, the production of two coloured final state particles, is one of the Standard Model processes with the largest cross-section. This fact makes quarks the optimal particles to probe for substructure in the early phases of the LHC experiment, since the low integrated luminosity is counteracted by their high production rates.

2.1.1 Contact Interactions

It is possible to stipulate many theories where quarks are composite objects made of preons. However, since no evidence is available on the possible nature of such a theory, any choice would be arbitrary, and would provide little to no insight. Instead, it is chosen to use an effective theory, where an extra four-Fermi contact interaction is added between four quarks. Such a model can only be used to describe the onset of quark compositeness; once the scale which governs compositeness is reached, the effective theory must be replaced with a more complete theory of the interaction between preons.

The simplest possible Lagrangian which couples four left-handed quarks is appended to the Standard Model QCD Lagrangian. Such a Lagrangian has traditionally been used
in past searches for compositeness [2, 3]:

\[ \mathcal{L}_{\text{compositeness}} = \mathcal{L}_{QCD} + \mathcal{L}_{qqq} \]

\[ \mathcal{L}_{qqq}(\Lambda) = \frac{g^2}{4\Lambda^2} \bar{\Psi}_q \gamma^\mu \Psi_q \bar{\Psi}_q \gamma^\mu \Psi_q \]

where \( \Psi^L_q \) are the quark fields, \( \Lambda \) is the compositeness scale, the interaction strength \( \frac{g^2}{4\pi} \) is arbitrarily chosen to be 1, and \( \xi \) is the parameter governing the interference with the QCD terms. The parameter \( \xi = +1 \) has been arbitrarily chosen to give destructive interference between the contact interaction terms and QCD. Previous studies using both types of interference demonstrated that this choice only changes the final limits on the compositeness scale \( \Lambda \) at the \( \sim 1\% \) level [4]. The added presence of a quark contact interaction to the QCD Lagrangian is expected to increase the dijet production rates as the kinematic regime approaches the compositeness scale \( \Lambda \).

2.2 Angular Distributions

In the well-known historical experiment performed by Geiger and Marsden, a gold foil was bombarded with alpha particles. By studying the resulting spread of the scattered particles, Rutherford inferred the substructure of the atom, leading to the formulation of the planetary model of the atom [5]. Analogously, it can be expected that the angular distributions of the outgoing jet pairs can potentially be probed to infer the presence of a substructure to the quark.

Standard Model QCD dijet production is dominated by t-channel exchanges, resulting in low angle scatters. However, the event topologies of contact interactions are expected to be more isotropic, leading to more jets being produced centrally. This analysis therefore consists of probing two dijet angular observables, the \( \chi \) distributions and the dijet centrality ratio \( (R_C) \), which have been shown in the past to exhibit sensitivity to the presence of quark contact interactions. If good agreement is found between the QCD predictions and the measured data, limits can be set on the existence of quark compos-
Chapter 2. Introduction

iteness up to a mass scale $\Lambda$, which can converted into a length scale at which the quarks are confirmed to be point-like.

2.3 Past Measurements

Various studies in the past have used dijet angular distributions as tests of QCD physics and to perform searches for quark compositeness. The first studies of dijet angular distributions were performed at the CERN SPS by the UA1 and UA2 collaborations. Both collaborations chose the center-of-mass scattering angle, $\theta^*$ as their observable, measuring the differential cross-section $\frac{d\sigma}{d\cos\theta^*}$ [6, 7].

The UA1 collaboration was the first to later use the $\chi$ distributions to probe the dijet angular distributions and set limits on quark contact interactions. The limits found by this search were set at $\Lambda > 0.42$ TeV [8]. Such studies were repeated at the Tevatron during both Run 1 and Run 2, gradually increasing the limits as more data was acquired [9, 10, 11]. The D0 collaboration was the first to use the $R_C$ observable, setting limits at $\Lambda > 2.4$ TeV [12]. Before the publication of the first LHC results, the highest limits had been set by the D0 collaboration, using the $\chi$ observable with an integrated luminosity of 0.7 fb$^{-1}$, at $\Lambda > 3.1$ TeV [4].

2.4 Overview

In this document, the methodology and results of a search for quark compositeness at ATLAS will be summarized. Chapter 3 describes the experimental apparatus used for this search, consisting of the accelerator complex (the LHC) and the particle detector (ATLAS). Chapter 4 focuses on the reconstruction of jets, the physics objects which are central to this analysis. In chapter 5, the angular observables used in this analysis are described, along with the event selection criteria. Chapter 6 offers a review of the Monte Carlo simulation techniques used to generate Standard Model predictions, and
summarizes the production of the samples. The main systematic uncertainties which affect the analysis are described in chapter 7. In chapter 8, the measurement is presented: a comparison is made between the measured data and the Standard Model Monte Carlo predictions, and limits are set on the existence on quark compositeness. Finally, chapter 9 provides a summary of the analysis and its results.
Chapter 3

ATLAS and the LHC

The Large Hadron Collider (LHC) at CERN is the world’s largest particle accelerator [13]. It spans a circumference of 26.7 km, and is located 150 meters underground, straddling the borders of France and Switzerland. The LHC was originally designed to collide two beams of protons at center of mass energies of 14 TeV, with instantaneous luminosities exceeding $10^{34} \text{ cm}^{-2}\text{s}^{-1}$.

The ATLAS (A Toroidal LHC ApparatuS) detector is one of the two general purpose detectors at the LHC [14]. Its purpose is to record the signature of the particles produced during the collisions of the two proton beams. Its geometry and operation are summarized in section 3.2.

3.1 The Accelerator

The LHC was built inside the 26.7 km long tunnel originally used by the LEP collider. The decision to use a tunnel which was designed for another experiment introduced design restrictions, but was preferred due to budget constraints. The physics requirements of the LHC program required a high instantaneous luminosity. The beam intensities required for such luminosities could not be achieved using anti-protons, leading to the decision to use two proton beams.
Since both beams use particles with the same charge, they must be subjected to two separate magnetic fields. Consequently, the beams must be carried in separate vessels. Due to space constraints in the tunnel, it was impossible to use two separate magnet systems. Therefore, a twin-bore magnet design was implemented, where a single cryostat accommodates both systems.

Since the length of the LEP tunnel could not be modified, larger bending magnetic fields are required to keep higher energy beams in the correct orbit. The magnetic fields required at the LHC are achieved by using superconducting magnets, operating at a temperature of less than 2 K, providing a design magnetic field of 8.33 T for 7 TeV beams. Radio Frequency cavities are used to accelerate the beams.

The injection chain of the LHC consists of a four stage process. The Linac2 first accelerates proton beams to 50 MeV. The Proton Synchotron Booster (PSB), Proton
Synchotron (PS), and Super Proton Synchotron (SPS) further accelerate the beam to energies of 1.4 GeV, 25 GeV, and 450 GeV, before the beam is finally injected into the LHC. Figure 3.1 shows a schematic view of the LHC complex and its injection chain.

### 3.1.1 Operating Conditions

During the 2010 data taking run, the LHC was operated at conditions below its original design, as part of the commissioning phase. The center of mass energy was lowered to 7 TeV, allowing the superconducting magnets to operate at lower currents, to prevent the possibility of a magnet system failure, such as the one which provoked the September 2008 accident [15].

The large instantaneous luminosities delivered by the LHC introduce important experimental concerns. The resulting presence of significant in-time and out-of-time pile-up must be accounted for during the experimental analysis of data. The in-time pile-up corresponds to the multiple interactions which are likely to occur during a single bunch crossing, due to the large number of protons in each bunch. The response of various detector electronics is characteristically longer than the spacing between each bunch; this leads to the presence of out-of-time pile-up. For example, at nominal run conditions, the full pulse shape in the liquid Argon calorimeters has a duration of more than 20 bunch crossings. Both phenomena (in and out-of-time) will produce energy readings which are not associated with the interaction which triggered the event being studied, leading to a degradation of the energy resolution and the introduction of possible biases.

### 3.2 The Detector

The ATLAS detector is sequentially divided into a series of subsystems. Each system is designed to exploit the specific energy loss mechanisms of different particles to discern their nature, and measure their kinematics. An overview of the ATLAS detector is shown
The inner detector, described in section 3.2.3, is closest to the beam pipe. It is used to reconstruct the tracks of charged particles. The magnetic field produced by the magnet systems (section 3.2.2) ensures that the tracks follow a curved trajectory, from which the momentum of charged particles can be induced.

The calorimeters (section 3.2.4) are used to stop particles which interact electromagnetically or strongly, and collect their energies.

Finally, the muon chambers (section 3.2.5) reconstruct the tracks of muons and other minimum ionizing particles which will pass through the calorimeters without depositing any significant energy. An extra magnetic field, provided by the toroidal magnet system of ATLAS (section 3.2.2), is used to obtain a momentum measurement.

The ATLAS detector is typically described as consisting of three fiducial regions, spanning different ranges of polar angles. The **barrel**, named aptly because of its cylindrical geometry, is situated in the most central region of the detector, and its subsystems
offer the best performance in terms of resolution and granularity. It is followed by the endcap region which has instrumental coverage from most subsystems, but typically with inferior resolution. Finally, additional detectors in the forward region complete the instrumentation, offering added detector hermiticity.

3.2.1 Coordinate Systems

Three different coordinate systems are used as a basis to describe the geometry of the detector. For each subsystem, the most natural coordinate system is then used.

In the Cartesian coordinate system, the z-axis is defined at the center of the beam-pipe, parallel to the direction of motion of the protons at the center of the detector. The y-axis points upwards (away from the center of the earth), and the x-axis points towards the center of the LHC ring.

The cylindrical coordinate system is better suited to the geometry of the tracking systems. The z-axis remains the same as in the Cartesian system. The variable $R$ represents the perpendicular distance to the origin of the z-axis, and $\phi$, the azimuthal angle in the (x,y) plane.

A spherical coordinate system, which follows the projective geometry of the calorimeters, is also used. In this coordinate system, the two angular variables, $\phi$ and $\theta$ represent the azimuthal and polar angles. Additionally, the pseudo-rapidity $\eta$ is often used instead of the polar angle $\theta$. It is defined as $\eta = -\ln(\tan(\frac{\theta}{2}))$. It is equal to the rapidity $y = \frac{1}{2} \ln(\frac{E+p_z}{E-p_z})$ for the case of massless particles.

Relative distances between objects are typically measured in $\Delta R$, which is defined as:

$$\Delta R = \sqrt{(\phi_1 - \phi_2)^2 + (\eta_1 - \eta_2)^2}$$

(3.1)

where $\phi_{1,2}$ and $\eta_{1,2}$ are the azimuthal angle and pseudo-rapidity of both objects.
3.2.2 Magnet Systems

The magnet systems in ATLAS provide large magnetic fields which serve to bend the trajectory of charged particles in the detector due to the Lorentz force. The curvature of the trajectory of particles can then be used in the tracking systems (the inner detector and muon spectrometer) to measure the momentum of particles. Proper modeling of the magnetic field generated in the tracking systems is crucial, as any uncertainty in the field strength will translate directly into an uncertainty on charged particle momentum determination. Figure 3.3 shows the geometry of the three magnet systems in ATLAS.

Muons in ATLAS will follow a unique trajectory due to the magnetic fields produced by the combination of the solenoid and toroid magnets. In the inner detector, the trajectory of a muon will curve along the azimuthal direction (due to the solenoid); in the muon chambers, it will curve along the polar direction (due to the toroid).
Barrel Solenoid

The magnetic field inside the inner detector is generated by the barrel solenoid. This solenoid, centered around the beam-pipe, generates a magnetic field of $\sim 2$ T at its center. Its walls are located at $R = 2.46$ and $2.56$ m away from the detector center. It extends up to $|z| = 2.9$ m from the interaction point, offering a full coverage of the inner detector. The field strength is very uniform across most of the useful volume enveloped by the solenoid; the field strength is constant within $\sim 5\%$ in the region $|z| \sim 1$ m, and gradually drops by up to $25\%$ at $|z| \sim 2$ m.

The location of the solenoid introduced an important design restriction. Because it is located in front of the EM calorimeters, the radiative length $X_0$ of the solenoid materials had to be kept to a minimum. This was done in order to reduce the probability of electromagnetic and hadronic particles interacting with the material before they reach the calorimeters. This led to a design which only contributes $0.66 \ X_0$ to the material in front of the calorimeters.

Barrel Toroid

The barrel toroid is one of the defining characteristics of the ATLAS detector. It consists of 8 large rectangular toroidal magnets symmetrically disposed in $\phi$. The magnets extend up to $|z| = 12.7$ m, and cover the space between $|R|$ of 9.4 and 20.1 m. The combination of these 8 magnets creates a cylindrical shell around the detector which is permeated by a toroidal magnetic field. The barrel muon tracking system is located inside this shell, ensuring that the path of charged particles in the tracker are deflected along the $\eta$ direction. Because of the complex geometry, the field is strongly dependent on position in the $(R, z)$ plane, with field strengths varying between 0.15 and 2.5 T.
Endcap Toroid

The endcap toroid follows the same principle as the barrel toroid. The 8 toroidal magnets are square in shape, and their field permeates the endcap muon systems. The magnets are located between $|R| = 1.7$ and $10.7$ m, and $|z| = 8$ to $12.5$ m. Similarly to the barrel toroid, the field is strongly dependent on position in the $(R,z)$ plane, with field strengths varying between 0.2 and 3.5 T.

3.2.3 Inner Detector

The inner detector consists of three subsystems, positioned with increasing distance from the interaction point. The spatial distribution of the energy deposited by particles traversing the inner detector are used to infer their trajectory. From these trajectories, the position of primary and secondary vertexes can also be interpolated. The momentum of the particles forming the tracks can also be measured. To ensure a good momentum resolution, the tracking system is therefore required to provide highly accurate spatial measurements. Figure 3.4 shows the lay-out of the three inner detector systems in both
the barrel and endcap regions.

**Pixel Detector**

The pixel detectors offer the finest spatial granularity, and are present in the region closest to the interaction point. This ensures that the best resolution is obtained closest to the beam, where the track density will be highest, and the vertexes will most likely be situated.

In the barrel, three cylindrical layers of pixel detectors are positioned at $R = 50.5$, $88.5$ and $133.5$ mm away from the beam center. These layers extend along the beam-pipe up to a distance of $|z| = 400.5$ mm. To extend the tracking coverage further up to the endcap, three additional layers, shaped as disks, are positioned at $|z| = 495$, $580$, $650$ mm away from the detector center. These disks offer instrumental coverage for the radius range of $88.8 < R < 149.6$ mm.

The layers are composed of 1744 separate sensors of dimensions $19 \times 63$ mm$^2$, each possessing 46080 individual discrete pixels of dimensions $50 \times 400$ µm$^2$, along with the associated control electronics. This yields a total of approximately 80 million read-out channels for the pixel detectors.

The sensor elements use silicon detector technology. The principle behind this technology relies on joining two silicon wafers which have been doped to create a diode-like structure [16]. The doping procedure introduces anomalous atoms in the crystal lattice of the silicon. These atoms, referred as donors (or acceptors), have more (or fewer) electrons available in the conduction (valence) energy band, changing the properties of the silicon lattice. The two different types of doping as referred to as p-type (using acceptor atoms) and n-type (using donor atoms). When a diode is created by joining a p-type and n-type wafer, a depletion region is formed at the interface of both materials. The extra electrons from the conduction band of the donors migrate to the holes in the valence band of the acceptors, creating a region with no mobile force carriers. An electric field
is generated across the junction by the ions created by the electron-hole migration. At steady-state, this field prevents any further electron-hole diffusion.

A reverse bias voltage is applied across the junctions, increasing the potential difference between the two regions, and widening the size of the depletion region. This also reduces the drift time of electrons, speeding up the response. When a particle excites the depletion region medium, electrons from the valence band migrate into the conduction band. The mobile charges then drift along the electric field, creating an electric current signal, which can be probed to detect the passage of particles in the sensitive element.

**SCT**

The SCT (SemiConductor Tracker) uses a similar technology and geometrical layout as the pixel detector. Four cylindrical layers are positioned in the barrel section at \( R = 299, 371, 443 \) and 514 mm away from the beam center. These layers extend along the beam-pipe up to a distance of \( |z| = 749 \) mm. The nine endcap layers are also disk-shaped, offering coverage between distances \( |z| \) of 854 mm and 2720 mm. The nominal disk coverage range is \( 275 < R < 560 \) mm. However, the coverage closer to the beam pipe diminishes as \( |z| \) increases, keeping the maximum polar angle coverage constant.

A total of 15912 sensors are present in the SCT detector, with each sensor holding 768 silicon micro-strips, for a total of over 10 million read-out channels. Each micro-strip is 12 cm in length, and 80 \( \mu \)m in pitch. To increase the precision of the position measurement, pairs of strips are superimposed, with a small angle (40 mrad) difference in their relative positioning. By combining the measurements from both strips, a more accurate position measurement can be made along the longitudinal direction.

**TRT**

The TRT (Transition Radiation Tracker) completes the inner detector array. Its barrel module extends from 554 to 1082 mm away from the beam in \( R \), with a coverage up to
$|z| = 712$ mm. The endcap module augments the coverage in the $848 < |z| < 2710$ mm and $644 < R < 1004$ mm region.

The TRT consists of cylindrical straws of 4 mm in diameter, installed parallel to the beam pipe. Because the straws are only oriented in one direction (parallel to the $z$ axis), the TRT provides only $R$ and $\phi$ measurements for each track hit. The straws are filled with Xenon based gas mixture, with a $31 \mu$m diameter anode located in the center of each straw. A typical track will traverse 36 straws as it moves past the inner detector. A large voltage is applied between the anode and cathode (the straw wall), to allow the collection of ionized electrons. The space between the straws is filled with fibers. The constant interface changes caused by the many fibers favors the emission of transition radiation between the straws.

The transition radiation photons, emitted by charged particles as they traverse materials with different dielectric constants, are absorbed in the gas mixture, generating larger signals than minimum ionizing particles. Because the emission rates of transition radiation are dependent on the particle’s Lorentz factor $\gamma = \frac{E}{m_c\gamma}$ (and hence, on its mass), the information from the TRT can be combined with the information from the calorimeters to improve the particle identification procedure.

While lacking the precision of the silicon based detectors, the larger spatial coverage and larger number of track hits the TRT provides ensures that it contributes significantly to the momentum measurement.

**Performance**

The ATLAS inner detector provides full tracking coverage up to $|\eta| < 2.5$. The coverage of the TRT is shorter (only up to $|\eta| < 2.0$), and therefore sets a limit on the electron identification range in the tracker. The inner detector was designed to provide a transverse momentum measurement resolution of $0.05\% \times p_T \oplus 1\%$.

It is important to note that the ATLAS inner detector is one of the most heavily
Figure 3.5: Geometry of all calorimeters in ATLAS.

instrumented tracking systems to date. The large density of material in the inner detector (varying between 0.5 and 1.1 $X_0$) means that a large fraction of electromagnetic particles will have already begun showering before reaching the calorimeters.

### 3.2.4 Calorimetry

The ATLAS calorimeter system is segmented into the barrel, endcap and forward regions of the ATLAS detector. The system is further segmented into electromagnetic (EM) and hadronic components, since the depths of the showers of EM and hadronic particles are governed by different characteristic length scales. The calorimeters offer a complete, symmetric coverage in $\phi$. One of the defining characteristics of the ATLAS calorimeters is their fine longitudinal segmentation. This segmentation allows the study of the development of EM and hadronic showers in terms of depth, allowing a better characterization of shower shape. This in turn can be used to improve the resolution of the energy measurements.

All calorimeter systems in ATLAS use sampling technology. Layers of active material,
capable of measuring energy deposits, are interspersed with layers of absorber material. The absorbers, typically made of high density materials, are only used to develop and stop electromagnetic and hadronic showers while offering no measurement of the actual energy deposits. Figure 3.5 shows the geometry of all calorimeters.

When describing the calorimeter, all granularity figures quoted are in units of $\Delta \eta \times \Delta \phi$.

**LAr EM Calorimeters**

The LAr EM calorimeters offer coverage in the range $|\eta| < 3.2$. For both the barrel and endcap sections of the calorimeter, the longitudinal layers closer to the inner detector have a finer spatial granularity, in order to obtain a better resolution in the position and shape measurement of the shower. The calorimeters use lead as the absorbing material, and liquid Argon (LAr) as their active material. An accordion geometry is used, ensuring an even detector coverage without any cracks in $\phi$. The calorimeters are designed to be projective in $\eta$, which means that the geometrical dimensions of the calorimeter cells vary to accommodate a constant $\Delta \eta$ granularity.

The barrel section covers up to $|\eta| < 1.475$, and is composed of 3 longitudinal layers with different read-out cell sizes. For most of the rapidity coverage, the first layer offers a granularity of $0.025/8 \times 0.1$, the second layer, $0.025 \times 0.025$, and the third layer $0.050 \times 0.025$. The endcap section offers coverage in the region $1.375 < |\eta| < 3.2$. It is segmented in three longitudinal layers at its center ($1.5 < |\eta| < 2.5$), and two layers at its two extremities. In the first layer, the granularity varies between $0.025/8 \times 0.1$ and $0.1 \times 0.1$, the latter occurring at the extremities. The second layer offers a granularity varying between $0.025 \times 0.25$ and $0.1 \times 0.1$, with the coarser granularity appearing again in the higher pseudo-rapidity sector. Finally, the third layer, only present at the center, has a granularity of $0.050 \times 0.025$. Figure 3.6 shows the three layers which compose a tower of the EM barrel calorimeter, displaying the accordion geometry of the electrodes and cell sizes.
Figure 3.6: Geometry of a tower in the EM Barrel calorimeter.

To improve the shower sampling of the calorimeters, a pre-sampler layer is added in front of both sections. The barrel solenoid contributes significantly to the amount of dead material located in front of the calorimeters. By having the pre-sampler located in front of the solenoid, a better measurement of the early development of the showers and energy loss due to the dead material can be made. The pre-sampler consists of active liquid argon layers (a single layer in the barrel, and two in the endcap). It offers a position granularity of $0.025 \times 0.1$.

The EM calorimeters offer excellent longitudinal shower containment. The material in the EM calorimeters adds up to a total of 25 to 40 $X_0$.

**Tile Calorimeter**

The Tile Calorimeter uses a combination of scintillator tiles as the active material, and steel as the absorber. This calorimeter offers coverage in the $|\eta| < 1.7$ range, and is separated into three modules, one barrel module and two extended barrel modules. All
modules are separated into three longitudinal layers; the read-out granularity is $0.1 \times 0.1$ for the first two layers, and $0.2 \times 0.1$ for the last layer. In the region between the barrel and the endcap, a special module named the tile gap, composed of only scintillator tiles, allows a partial measurement of the energy of particles falling in this detector crack. The total material in the tile calorimeters adds up to a total of approximately 7.4 interaction lengths ($\lambda$).

**Hadronic Endcap**

The Hadronic Endcap calorimeter (HEC) covers the $1.5 < |\eta| < 3.2$ pseudo-rapidity range, and uses a combination of liquid Argon for the active material, and copper for the absorber. The two HEC modules are disk shaped, and are divided into two longitudinal layers, for a total of four layers. The front module covers the $1.5 < |\eta| < 3.2$ pseudo-rapidity region, while the the back module offers coverage in $2.5 < |\eta| < 3.2$. The granularity in all layers is $0.1 \times 0.1$ in the front module, and $0.2 \times 0.2$ for the back module. The hadronic endcap provides a total of $\sim 10 \lambda$.

**FCAL**

The FCAL extends from $3.1 < |\eta| < 4.9$. The large coverage of the FCAL in pseudo-rapidity contributes to the excellent hermiticity of the ATLAS detector. Its three layers provide electromagnetic and hadronic energy measurements. The active material is liquid Argon, and the absorber is copper for the first layer, and tungsten for the second and third layers. The geometry of the FCAL is different from the other calorimeters: electrode rods reside in tubes filled with liquid argon, which are positioned parallel to the beam pipe, inside the absorber matrix. The FCAL offers a total of approximately $10 \lambda$ interaction lengths.
Performance

The calorimeters were designed to achieve a good resolution in the energy measurement of electromagnetic and hadronic particles, as defined by the physics requirements of the LHC program. For electromagnetic particles, an energy resolution \( \sigma_E \) of \( \frac{0.2}{E(GeV)} \oplus \frac{10\%}{\sqrt{E(GeV)}} \oplus 0.7\% \) is expected, in the barrel and endcap regions. For hadronic particles, the resolution is much larger, due to the intrinsic fluctuations in the hadronic showers. A resolution of \( \frac{50\%}{\sqrt{E(GeV)}} \oplus 3\% \) is expected for the barrel and endcap, and a resolution of \( \frac{100\%}{\sqrt{E(GeV)}} \oplus 10\% \) is expected for the forward region. Early studies using in-situ methods show good agreement between the predicted and measured resolutions, for the barrel and endcap regions. Figure 3.7 shows the number of interaction lengths as a function of \( \eta \), to demonstrate the good shower containment of the ATLAS calorimeters.

3.2.5 Muon Chambers

Due to their long lifetimes, muons do not decay inside the detector. Ionization is the primary energy loss mechanism of muons at the energy range they are typically produced at. 
Mechanisms with larger $\frac{dE}{dx}$, such as bremsstrahlung, have low associated cross-sections. Therefore, muons behave largely as minimum ionizing particles, and will traverse the entire detector (including the calorimeters) with minimal energy losses. To identify muons and measure their momentum, four tracking subsystems are used, positioned outside the calorimeters. The systems use the deflection of the muons tracks, due to the various magnetic fields permeating the muon chambers, to measure momentum.

The muon chambers are separated into two categories. The precision tracking chambers offer high granularity position measurements, while the trigger chambers offer very fast, but coarse, track identification, which can be used for event triggering. Figure 3.8 shows the location of all muon chambers.

**MDT**

The Monitored Drift Tubes are the main component of the precision tracking chambers. The chambers cover the range up to $|\eta| < 2.7$. In the barrel, a series of 3 chambers are located at $R = 5$, 7.5 and 10 m. Each chamber contains 6 to 8 tubes layers. To provide
full coverage in $\phi$, 16 azimuthal sections are used. These sections achieve a $\phi$ symmetry; small sections are positioned as the faces of an octagon, centered on each of the toroids, and large sections complete the coverage between each toroid pair. In the endcap, 4 chambers are positioned at $|z| = 7.4, 10.8, 14$ and 21.5 m away from the detector center. This ensures that, on average, a muon track will cross 3 different chambers. Each tube uses a single electrode, and uses the signal generated by the drift of electrons created during the electron avalanche.

Due to their layout, being placed orthogonal to the beam-pipe, the MDTs do not offer any precise measurement in the $\phi$ plane. Therefore, multiple tracks in the same detector region will create an ambiguity in matching the various hits. However the probability of finding such closely positioned tracks has been determined to be negligible. In the rare cases of highly boosted particle decays to two minimum ionizing particles, the inner detector can be used to aid in track identification.

**CSC**

The high rapidity tracking region, $2.0 < 4|\eta| < 2.7$, has the largest expected density of tracks. In this busy environment, the probability of finding multiple tracks in the same region is much higher. The Cathode Strip Chambers address this problem by combining measurements from orthogonally positioned read-out strips in multi-wire proportional chambers. A single CSC chamber is used to replace the first MDT layer in the high rapidity tracking region.

**RPC**

In the barrel region ($|\eta| < 1.05$), Resistive Plate Chambers are used as the trigger chambers. The RPCs are composed of two separate units. Each unit uses two read-out strips, one on each side of the gap, which are positioned orthogonally. This provides a position measurement in both the $\phi$ and $\eta$ plane. The units are wireless gas tracking chambers:
a pair of resistive plates surround a gas gap with a large, uniform electric field created by two parallel high voltage plates. As an electron avalanche is triggered in the gas due to the passage of an ionizing particle, the signal is read out by the strips via capacitive coupling.

Two RPC layers are positioned at the inner and outer face of the second MDT chamber, and the third RPC layer is located either at the inner or outer face of the third MDT chamber, respective of whether it belongs to a small or large section.

**TGC**

In the endcap region ($1.05 < |\eta| < 2.4$), Thin Gap Chambers are used as the trigger chambers. These chambers use the same principle as multi-wire chambers. However, the distances between the components of the TGCs are very small: the anode wires are 1.8 mm apart, and the distance between the anodes and the cathode layer is 1.4 mm. The combination of the large high voltage and the compact geometry ensures that the response time is very fast. The TGCs provide an azimuthal measurement due to the orthogonal positioning of the cathode strips with respect to the wires.

Seven layers of TGCs, split into two doublets and a triplet, are used to complement the middle MDT layer in the endcap. For the inner layer, two sets of single TGCs are used to cover the region in front of the MDTs.

**Performance**

The muon system offers track momentum measurements up to $|\eta| < 2.7$, while the triggering is limited to $|\eta| < 2.4$. The muon system provides precision position measurements along the direction where tracks are expected to bend (the polar axis). Azimuthal information can also be obtained by combining the information from track hits in the trigger chamber, which offer a modest spatial resolution. The spatial resolution of the trigger chambers is sufficient to provide a momentum measurement which can be used for trig-
ger thresholds. The combined information from the muon spectrometer is expected to provide a transverse momentum resolution of $10\% \times p_T$ for tracks of up to 1 TeV.

### 3.2.6 Trigger

At design luminosities, collisions will occur at a rate of approximately 1 GHz. Unfortunately, the data acquisition systems, due to hardware limitations, can only accommodate the bandwidth (rate of data transfer) associated with an event rate of up to $\sim 400$ Hz.

The ATLAS detector uses a three-level trigger system to limit the rate at which events are recorded, in order not to exceed the maximum supported bandwidth. The three levels are designated Level 1 (L1), Level 2 (L2) and Event Filter (EF), and will be described in the following sections. During the data runs used for this analysis, the L1 trigger was used exclusively to limit the event rate to the desired bandwidth. The remaining triggers (L2 and EF) were in pass-through mode.

#### L1

The L1 trigger uses information from the calorimeter and muon systems to coarsely identify the signatures of different physics objects. This information can then be used to accept or reject events based on simple object-based criteria, in order to reduce the acceptance rate to a design value of 75 kHz. This system is implemented at the detector level using a set of dedicated fast electronics. The system is designed to provide a decision within $2.5 \, \mu s$ of the bunch-crossing time. Because the collision rate is much faster than this response time, a pipeline is used in the L1 systems to store the events until the decision is made. The L1 systems return a set of parameters to the Central Trigger Processor (CTP). These parameters can be used as criteria for the various trigger selection menus at L1. The information available consists of the numbers of each object observed, and the thresholds which they satisfy.

The muon system uses a simple track identification approach. One of the middle
layers in the RPC/TGC is used as a pivot point for track fitting. The track fitting procedure consists of defining a set of 6 different roads, corresponding to the expected track curvature, starting from the interaction point, for a predetermined set of $p_T$ ranges. The roads define a spatial range in which a track hit is to be expected in each layer. Each hypothesis is tested, and if a sufficient number of coincidences are found, a muon candidate is passed to the CTP along with the $p_T$ threshold it satisfied. The coincidence requirement greatly reduces the contributions from noise-induced fake tracks and cavern backgrounds.

In the calorimeter, the L1 system allows the selection of events containing electrons/photons, jets and tau leptons decaying hadronically. Also, events can be selected based on missing transverse energy and transverse energy scalar sums. The calorimeter is coarsely segmented at L1 into trigger towers of size $0.1 \times 0.1$ (in $\Delta\eta \times \Delta\phi$) from which the energy contributions from the electromagnetic and hadronic calorimeters are read out. The electron/photon and tau lepton triggering extends up to $|\eta| < 2.5$, the jet triggering, up to $|\eta| < 3.2$, and the missing transverse energy and transverse energy scalar sums use the information from the region $|\eta| < 4.9$, to ensure a more hermetic coverage.

Jets are built using square combinations of jet elements ($2 \times 2$ trigger towers), forming areas of 0.4, 0.6 or 0.8 in $\Delta\eta \times \Delta\phi$. The $E_T$ sums for all these possible sliding windows are compared to pre-determined thresholds to decide whether a jet candidate is reconstructed. To avoid overlap in the jets, each jet candidate is required to be centered on a jet element corresponding to a local $E_T$ maximum. Electron and photon identification is performed using $2 \times 1$ combinations of trigger towers. Further selection criteria are based on the isolation (energy in nearby trigger towers), and hadronic leakage (energy in the hadronic calorimeters). The tau algorithm, similarly, uses the information from both electromagnetic and hadronic calorimeters, combined with isolation veto criteria.
L2

Once an event has satisfied one of the L1 trigger criterion, based on combined multiplicity and threshold information from the muon and calorimeter systems, the decision is sent to the read-out drivers (RODs). The detector region which contains the trigger object is referred to as the Region of Interest (RoI). The complete detector information associated with the RoIs is sent from the RODs to the L2 system. Focusing only on the RoIs, which only contain < 2% of the total event information, limits the amount of data which has to be transferred and processed by the L2 system. In addition to the information from the muon and calorimeter systems, the L2 system also uses information from the inner detector.

The L2 system uses a farm of computer processing units to perform trigger decisions. A series of refined criteria are used to determine whether the event satisfies any of the possible pre-defined L2 signatures. Due to the longer allowed processing time of \( \sim 40 \) ms, a more complex reconstruction procedure can be applied to the trigger objects in the RoIs. The L2 system is designed to reduce the acceptance down to a rate of 3.5 kHz.

EF

Similarly to the L2 system, the EF uses a computing farm to perform event selection. At the EF level, a full event reconstruction is performed, using software analogous to the ATLAS offline analysis software, using information from the full detector. Once an event has passed the EF, it is sent to the CERN data recording facility, and stored to disk. The events which pass the EF are classified according to which ATLAS physics stream they belong to. These streams in ATLAS are: electrons, muons, photons, tau leptons, missing transverse energy, jets, and B-physics. It is important to note that one event can be classified into more than one stream, and will therefore be copied to the corresponding number of output files. The refined selection criteria used in the EF ensures the event acceptance rate is reduced down to the required 400 Hz.
Chapter 4

Jet Reconstruction

The partons present in the final state of a fixed order perturbative QCD calculation cannot be directly observed in a detector. Before reaching the detector, the partons undergo a complex series of interactions (fragmentation), usually factorized into two parts: showering and hadronization. During the showering stage, extra partons are radiated from the hard scatter final-state partons. Then, during the hadronization stage, these partons are joined into colour singlet bound states (hadrons). The large showers of particles reaching the detector are grouped into jets: objects which represent an approximation of the kinematics of the partons present in the final state of the hard scatter. This approximation is loosely defined, as the separation of the hard scatter and the subsequent parton showers into different processes is simply an approximation used in the Monte Carlo pertubative calculations.

The jet topology of an event is entirely determined by the algorithm, called a jet finder, which dictates how the input constituents will be joined, and how the kinematics of the output jets will be constructed. Because jet finders simply recombine constituents, they can be used on any form of input with well-defined kinematics: partons, hadrons, and even detector-defined quantities (inner detector tracks, calorimeter segments, etc.).
4.1 Jet Finder Algorithms

The typical jet algorithms use a measure of proximity between input constituents to join the constituents into single jets. Most algorithms require the definition of a typical scale, often called the distance parameter ($R$). This scale defines the distance over which constituents will be separated into two individual jets. This choice of scale is entirely arbitrary, but is usually chosen such that the distance defined characterizes the typical spatial extent of the soft radiation from a single parton. This ensures that most of the energy of a final-state parton would be included in a single jet. Such a definition is signature-dependent.

Two recombination schemes are available to calculate the final jet kinematics. The original scheme, labeled Snowmass, consists of weighting the constituent kinematics in detector space by their transverse energy, as illustrated [17]:

\[
E_T^{\text{jet}} = \sum E_T^{\text{const.}},
\]
\[
\eta^{\text{jet}} = \frac{1}{E_T} \sum E_T^{\text{const.}} \eta^{\text{const.}},
\]
\[
\phi^{\text{jet}} = \frac{1}{E_T} \sum E_T^{\text{const.}} \phi^{\text{const.}}.
\]

(4.1)

where $E_T^{\text{const.}}$, $\eta^{\text{const.}}$, and $\phi^{\text{const.}}$ are the transverse energy and detector coordinates of the input constituents. The mass of the jet can be approximated after calculating the transverse momentum of the jet:

\[
p^{\text{jet}}_{x,y} = \sum p^{\text{const.}}_{x,y},
\]
\[
p^{\text{jet}} = \sqrt{(p^{\text{jet}}_x)^2 + (p^{\text{jet}}_y)^2}
\]

(4.2)

This leads to an incorrect approximation of the mass in the cases where a large separation between the constituents occurs in $\phi$ space, due to the definition of $E_T$.

The more recently used scheme is the four-vector recombination scheme. In this case, the constituents are assumed to be massless, compared to the scale of the jet, and their kinematics are added using a vector sum in momentum space:

\[
(p_x, p_y, p_z, E)^{\text{jet}} = \sum (p_x, p_y, p_z, E)^{\text{const.}}.
\]

(4.3)
In this case, the jet gains mass as constituents are merged. However, correctly measuring the mass of a jet requires that single hadrons be used as constituents.

The Snowmass scheme does however have advantages of its own, due to its simpler definition. This algorithm is much easier to implement computationally. Also, since $E_T$ is defined as a scalar sum, these vectors transform much more simply under Lorentz boosts. Furthermore, in the presence of negative energy constituents (due for example to detector noise), the Snowmass scheme can still offer a well-defined measure for jet kinematics.

Various considerations must be taken into account when choosing which jet finding algorithm is to be used. Some of these considerations, and possible choices for jet algorithms, are detailed in the next sections.

### 4.1.1 Requirements

Jet algorithms are used to define a common ground for comparison between predictions obtained using fixed-order perturbative QCD Monte Carlo calculations, and the hadronic final states reconstructed in the detector. Therefore, an ideal jet algorithm is one which will return identical (or similar) jet final states independent of the choice of the type of input constituents (partons, hadrons, reconstructed detector objects, etc.). This is the primary consideration for the quality of a jet algorithm [18]. One of the important requirements of jet algorithms is therefore infrared safety. An infrared-safe jet algorithm is one that will return the same jet topology independently of the presence of soft emissions in the final state.

A good jet algorithm must also be simple to implement experimentally; for modern particle physics experiments, where large datasets must be analyzed quickly and efficiently, this means that the algorithm must have a good computational performance. Also, an algorithm which returns a well-defined and regular jet area is highly preferable. This is due to the fact that the energy contributions to jets from sources other than
the hard scatter (multiple interactions, detector pile-up, initial state radiation, etc.), are easier to characterize for jets with well-defined areas. Finally, it is important that the jet algorithm be robust under pile-up conditions (see section 3.1.1). This consideration is slightly different from the requirement of infrared safety, because in this case, the extra (soft) energy emissions come from a different source than the original hard scatter under study. It is therefore important that the jet topology not be significantly affected by the presence of pile-up.

4.1.2 Cone Jets

Because of the circular nature of hadronic showers in the $\eta$ and $\phi$ representation in which the detector is built, a natural choice is to attempt to cluster the shower particles using a projective cone, starting from the collision vertex. This ensures that the input constituents are joined based on their geometrical proximity. The traditional method consists of seeding cones around input jet constituents, provided they satisfy an energy significance requirement. The transverse-energy-weighted centroid of the cone is calculated, and if it differs from the geometrical center, the cone is re-located at its centroid. This procedure is iterated until the cone is found to be stable. Once all possible cones have been found, there is a possibility that some of the constituents may be included in two different cones, in which case the cones are considered as overlapping. This problem is resolved, to avoid the double counting of energy, by the \textit{split/merge} process. If more than a fraction of the jet’s energies is shared, the jets are merged together. Otherwise, they are separated, and the constituents are distributed amongst the two cones following a pre-determined procedure.

The ATLAS cone algorithm places an energy requirement of 1.0 GeV on its seed constituents, and uses the four-vector recombination scheme to calculate the jet kinematics [19]. Once a constituent has been used as a seed, it is removed from the list of possible seeds. This procedure leads to the problem of \textit{dark towers}: if the cone migrates
far enough from its original seed such that the seed is no longer included in the cone, there is a probability that no other jets will include this seed [20]. This can lead to situations where constituents representing significant energy deposits are absent from the jet final state. This pathological behavior is one of the well-understood flaws in the cone algorithm.

The ATLAS cone algorithm also displays problems with infrared safety. Two high energy input constituents at a distance larger than $R$ but smaller than $2R$ of each other can either be conglomerated into a single jet, or be separated into two jets. The presence of soft radiation between the two constituents can ‘pull’ a cone seeded by one of the two constituents, during the iterative process, towards the other constituent. This is another major flaw on the cone algorithm.

The ATLAS cone algorithm uses a split/merge energy fraction of 0.5. If two jets are split, the constituents are distributed based on their angular proximity in $\eta$ and $\phi$ space to the centroid of each cone. Due to the split/merge procedure implemented, the jet areas produced by the cone algorithm are not very regular; the merging of two jets leads to a very un-cone-like jet shape.

Various attempts have been made at fixing the pathologies inherent to the cone algorithm. The midpoint algorithm, for example, attempts to solve the $R$-$2R$ problem by trying to seed a cone at the middle point between every pair of cones [21]. The problem of dark towers can also be solved; for example, the $JETCLU$ algorithm used at CDF used a ratcheting procedure which ensured that the seed towers were always kept inside the final jet area [22].

The seedless infrared safe cone algorithm (SISCone) was also investigated at ATLAS [23]. By removing the concept of seeds from the cone algorithm, it becomes infrared safe. However, the performance of the algorithm suffers greatly due to the computational complexity: with no seeds to give anchoring points for the cones, a large number of cone permutations must be calculated.
4.1.3 Recombination Algorithms

The recombination algorithms offer an alternative approach inspired by theoretical considerations. In these algorithms, pairs of input constituents within a pre-determined minimal distance to each other are merged into a new constituent (often called ‘protojet’). This procedure is repeated until no more merging is possible, at which point an approximation to the partonic final state is fully determined. This approach can be seen as analogous to the reverse of the fragmentation associated to partons; input constituents are rejoined, reversing the effects of the particle ‘splitting’ due to radiation, until the initial partons are reconstructed from them.

The traditional clustering algorithm uses the relative $k_T$ of input constituents as a metric [24]. The distance between constituents is therefore calculated as:

$$d_{i,j} = \min(k_{T_i}^2, k_{T_j}^2) \frac{\Delta y^2 + \Delta \phi^2}{R^2}$$

(4.4)

where $\Delta y$ and $\Delta \phi$ are the rapidity and azimuthal difference between two input constituents, and $k_T$ is the transverse momentum of each input constituent. The distance between a constituent and the beam is defined as:

$$d_{i,B} = k_{T_i}^2$$

(4.5)

and constituent pairs obeying:

$$d_{i,j} < d_{i,B}$$

(4.6)

are merged. As is evidenced by the definition of $d_{i,j}$, this metric ensures that lower transverse momentum constituents are preferably merged with the closest large transverse momentum constituents. This is consistent with an approximative picture of jet fragmentation, where large-angle emissions have low transverse momenta relative to the parton which spawned them. This behavior, however, leads to difficulties with robustness under pile-up conditions. The jet topology is heavily driven by the low $k_T$ components, and the pile-up contributions introduce many such components. Therefore, under heavy pile-up environments, the $k_T$ algorithm may return different final jet topologies.
However, the $k_T$ metric is not unique: other metrics can be produced, by varying the original expression. Two other alternatives can be obtained by changing the exponent used. The algorithm using the distance parameters:

$$d_{i,j} = \min(k_{Ti}, k_{Tj}) \frac{\Delta y^2 + \Delta \phi^2}{R^2}$$

is referred to as the Cambridge-Aachen algorithm. In this algorithm, the sequential recombination only uses information from the angular separation of constituents. The advantage of this algorithm is the hierarchical angular substructure it creates. If the merging steps are undone, the substructure of the jet can be easily probed, as the last merging is always the one with the largest angular separation. This algorithm is therefore particularly advantageous when attempting to probe the substructure of jets, especially in the cases where highly boosted objects are formed into single jets.

Of particular interest is the variation with a negative exponent:

$$d_{i,j} = \min(k_{Ti}^{-2}, k_{Tj}^{-2}) \frac{\Delta y^2 + \Delta \phi^2}{R^2}$$

referred to as the anti-$k_T$ algorithm [25]. This algorithm behaves in an opposite manner to the original $k_T$ algorithm, in the sense that high transverse momentum constituents are merged first. This behavior produces many desirable features. In the absence of further hard constituents, the algorithm will merge all low transverse momentum constituents in near perfect cones around the hard constituents. This yields jets with very regular areas. Also, this algorithm is pile-up safe, as jets are preferentially centered around the hard constituents. However, unlike the $k_T$ algorithm (which is hierarchical in relative $k_T$) and the Cambridge-Aachen algorithm (which is hierarchical in angle), the anti-$k_T$ algorithm offers no insight on the jet substructure. Still, due to the many advantages it offers over the other algorithms, it has been chosen as the preferred default at ATLAS.

Figure 4.1(a) demonstrates the very regular jet areas created by the anti-$k_T$ algorithm,
compared to the $k_T$ algorithm (figure 4.1(b)) [25]. A grid of soft ghost particles was generated on top of a parton-level simulated event, for the sake of illustration.

4.2 Jet Inputs

4.2.1 Detector Towers

Because it is impossible to directly reconstruct individual hadrons, the jet finding must be performed on a different type of input. The traditional approximation is to use the four-vector of each detector tower (0.1 by 0.1 in $\eta$ and $\phi$ space) as an input particle. Each tower is assumed to be massless, since the mass of hadrons is much smaller than the characteristic energy of the jets studied. This tower size was found in the past to be a reasonable approximation for the shower size of a single hadron. The resulting jets are known as TowerJets.

There are difficulties associated with this method, however. Because the noise profile of Liquid Argon detector cells is quasi-symmetric around zero [26], it is not unlikely for towers to have a negative total energy. This becomes problematic during the jet finding procedure, since negative energy four-vector algebra is ill-defined. Various methods have
been attempted to address this problem. Using only positive energy towers introduces a significant bias to the energy of the jets. This results in spurious jets spawned by noise in the calorimeter. One possible solution is to find a method to keep the information about the negative noise component, while removing all negative energy input vectors.

The ‘TowerNoise’ approach, which used to be the standard at ATLAS, attempts to resolve this problem. A grid of 7 by 7 towers (0.7 by 0.7 in $\eta$ and $\phi$ space) is constructed around each negative energy tower. Positive energy towers in this grid are merged with the seed negative energy tower, until the resulting energy is positive, starting with the nearest, lowest energy towers. While mostly addressing the problem of energy biases, this approach introduces another pathological behavior of its own. Because of the large geometrical reach of the grid, needed for cancellations to take place, non-local effects can occur; in extreme cases, a resulting merged tower can have constituents as far as $\sim 0.5$ in $\eta$ and $\phi$ space from its centroid. Since the jet finding is performed on these merged towers, this can lead to very irregular jet shapes.

The newest proposed alternative is the GhostTower scheme. In this scheme, negative energy towers are replaced by ghost towers, with the same geometrical position, but a negligible energy (0.001 MeV), during the jet finding procedure. Because of their low energy, these ghosts towers are effectively akin to ‘very soft radiation’ and do not have any pathological effects on infrared-safe jet finding algorithms. During the jet recombination, the resulting four-vector from the positive energy and ghost towers is re-scaled by the full energy contained in the original towers. This ensures that the energy cancellations due to the quasi-symmetric noise profile still take place, but that the negative energy towers do not affect the direction of the jet. The ‘GhostTower’ approach was found to perform much better than the ‘TowerNoise’ approach; while exhibiting the same desired noise cancellation, ‘GhostTower’ jets did not introduce any pathological behavior.
4.2.2 Topoclusters

The TowerJet approach highlighted above was shown to have many limitations. The coarse tower granularity, motivated by detector cell sizes and estimates of single particle shower sizes, does not make full use of the excellent ATLAS calorimeter granularity in the innermost layers. Also, many problems arise due to the presence of noise in the calorimeter.

The ATLAS topological clustering algorithm was designed to address these limitations [27]. This algorithm attempts to reconstruct the signature of single particles in the detector, while implementing a noise suppression algorithm. This is done by grouping the energy deposits in cells using a three-dimensional nearest neighbor approach, as described below:

- A cluster is seeded by the largest energy unclustered cell with \(|E_{cell}| > 4 \sigma_{noise}\), where \(\sigma_{noise}\) is the characteristic electronic noise for a cell, measured during dedicated runs

- The cluster is iteratively expanded by adding every cell with \(|E_{cell}| > 2 \sigma_{noise}\) geometrically adjacent to the cluster

- Once no more cells with \(|E_{cell}| > 2 \sigma_{noise}\) are found in contact with the cluster, the last outside layer of cells bordering the cluster is added to it

Any cell already included in a cluster is removed from the list of cells available. Once all possible clusters have been formed, a splitting algorithm is applied to the clusters. This algorithm splits a cluster if multiple local energy maxima are located inside. The splitting step was designed to isolate the signature of single particles inside individual clusters. However, performance studies in simulations demonstrate that a single cluster may contain the energy signature of many particles.

Topoclusters are the favored jet inputs at ATLAS, due to the advantages mentioned above. However, study of the tower jets is also very important as a cross-check, since
they are much simpler constructs. Topoclusters are very sensitive to the noise modeling in the calorimeter. Incorrect noise thresholds could lead to pathological cluster formation, leading to potentially extreme cases, such as, for example, events with a single cluster spanning the entire detector. However, constant online monitoring of the jets and their constituents, with comparisons between tower and topocluster jets, ensures that no problems arise.

4.3 Jet Energy Calibration

Because of the non-compensating nature of the ATLAS calorimeter, the calorimeter response of hadrons is lower than the response of electromagnetic particles. This motivates the need for dedicated calibrations to be applied to jets. Jets can be calibrated at two different levels:

- **Particle level**: Jets are calibrated to reflect the energy of the particles contained in the jet area. This calibration effectively only addresses the effects of the lower calorimeter response to hadrons.

- **Parton level**: Jets are calibrated to reflect the energy of the parton which spawned them. This includes more effects than the particle level calibration. For example, the effects of hard final state radiation causing energy to fall outside of the jet area (out-of-cone) must be, on average, corrected for. This approach is particularly useful when making comparisons between data and fixed order theory predictions, where radiation is not taken into account by the predictions.

The official calibrations in ATLAS are provided at the particle level. This is motivated by the complications introduced in attempting to provide parton level corrections. The parton level corrections are very signature-dependent, and a different set of calibration constants would therefore be needed for every specific final state.
The lower calorimeter response to hadrons is due to various effects, and the particle level calibration scheme attempts to correct for the following effects:

- **Non-compensation**: The response of hadronic showers is lower than the response of electromagnetic showers. Electromagnetic showers are quite predictable: the energy of incoming photons or electrons is eventually used to ionize the detector medium. However, hadronic showers are considerably more complicated due to the advent of strong interactions which can occur between the hadrons and nuclei in the detector. These strong interactions can result in the production of new particles which change the nature of the shower. Some of the energy in the shower can therefore be lost due to the creation of weakly-interacting particles, or simply absorbed by the media through nuclear interactions.

- **Dead material**: The incoming hadrons comprised in the jets can lose some energy by interacting with the inactive material in the detector.

- **Shower Containment**: The ATLAS calorimeters have been designed to optimize the longitudinal containment of the hadronic showers. However, it is still possible for some of the energy of particles in the shower to leak outside the detector.

- **Detector Out-of-cone**: Detector out-of-cone effects are corrected for in the particle level scheme, and represent the cases where a hadron is inside the jet area, but some of the particles created during its interaction with the detector fall outside of the jet area.

Various calibration schemes have been considered in ATLAS; they will be summarized in the following sections.
4.3.1 EMJES

The only officially supported scheme for early data is the so-called ‘EMJES’ (Jet Energy Scale correction based on the ElectroMagnetic jet kinematics). In this scheme, jets are corrected based on their position in the calorimeter ($\eta$), and transverse momentum ($p_T$). The correction is applied using a simple scale factor which scales the momentum and energy of the jet equally. This scheme therefore assumes that the relationship between momentum, energy and mass of the jet is correct at the electromagnetic scale. This calibration is highly dependent on the jet momentum: as the jet energy increases, the fraction of the jet’s energy deposited through electromagnetic showers increases. This is due to the fact that $\pi^0$ mesons decay preferentially into 2 photons. Jets consist mostly of $\pi^{0,\pm}$ and $K^{\pm}$, and the charged hadrons have a probability of producing a $\pi^0$ meson when interacting with the detector. Therefore, the more interactions occur in the hadronic shower, the more likely the $\pi^0$ mesons will constitute a large part of the shower. Therefore, very high energy jets require very little compensation, as most of their energy is deposited through the electromagnetic showers of photons produced by $\pi^0$ decays.

The calibration constants are derived by studying the detector response of jets in Monte Carlo samples. The samples used are the same as used for the analysis, and are described in section 6.4. The full ATLAS detector simulation, further detailed in section 6.2, simulates the passage and interactions of particles through the detector. Jets built from the Monte Carlo event generator particles (Truth jets) are matched to the jets built from the simulated calorimeter energy deposits (Reco jets). By comparing the momenta of the matched jets, an average correction is derived as a function of $p_T^{truth}$, for a variety of $\eta^{truth}$ detector slices. This function is then re-mapped as a function of $p_T^{reco}$; this procedure is called numerical inversion. The response curves for each $\eta$ detector slice are fitted with the following functional in an effort to reduce the effects of statistical
fluctuations:

\[ f(p_T) = a_0 + \sum_{i=1}^{4} \frac{a_i}{\log(p_T)^i} \] (4.9)

The selection and matching criteria used by the method are worthy of note, as they define the conditions under which the calibration is expected to provide the right correction. A pair of Truth and Reco jets is only matched if both are within a detector distance \( \Delta R \) of 0.3. Also, only isolated jets are matched; a jet is considered isolated if no other jet is present within a distance \( \Delta R \) of 2.5 \( R_{\text{cone}} \). Non-isolated jets are known to have a lower response, and would contribute to a lower response tail which would complicate the fitting procedure used.

It is important to note that this correction only attempts to correct the average momentum of jets; no attempt is made at improving the momentum resolution. The uncertainty associated with this choice of calibration is described in details in section 7.1. Figure 4.2 shows the calibration constants as a function of jet \( p_T \) for two different detector \( \eta \) slices.
4.3.2 Global Cell Weighting

The ‘GCW’ scheme (Global Cell Weighting) makes use of the fine longitudinal segmentation of the calorimeters to improve the resolution of the jet energy measurement. Each cell in a jet is assigned a particular weight based on the energy density \( \frac{\text{Energy}}{\text{Volume}} \) in the cell, and the calorimeter sampling the cell belongs to. The weights are derived using a Monte Carlo study, comparing Truth Jets and Reco Jets, similar to the method highlighted above. An optimization procedure is used to obtain the set of weights which minimizes the jet energy resolution. This method was originally designed by the H1 collaboration [28].

Because the optimization is made on the resolution, it is not guaranteed that the absolute scale will be restored by the correction. Therefore, the GCW calibration also has to be followed by another correction identical to the ‘EMJES’, with constants re-derived for GCW calibrated jets.

This scheme also allows for corrections of the jet direction. Jets can span multiple detector regions employing different technologies and therefore possessing different responses to hadrons. This can lead to a bias in the jet direction towards the region with the higher response; the individual cell weights will offer greater compensation in the region with lower response, yielding a better measurement of the jet position.

The ‘GCW’ approach is still undergoing validation using collision data. This is due to the added complexity associated with validating such a scheme; while the EMJES scheme only depends on the correct description of the response of hadrons in the calorimeter simulation, the GCW has an added dependence on the shape of the jet shower. Different shower profiles will result in different distributions of the shower energy in the various calorimeter segments. For this reason, this scheme is not yet used in ATLAS analyses.
4.3.3 Local Hadronic Calibration

The ‘LC’ (Local Hadronic Calibration) scheme uses the capacity of the topoclustering algorithm to reconstruct the complex three-dimensional signature of calorimeter showers from single particles to offer a more sophisticated calibration scheme [29]. The calibration is applied at the cluster level, and is aimed to correct for detector non-compensation and energy losses in dead material.

In this scheme, each cluster is classified as containing either an electromagnetic or hadronic shower. This classification is based on cluster ‘moments’: variables which describe the shape of the energy deposit, using information about the shower depth, width, energy density, etc. The classification is based on probability weights derived from the distributions of the chosen variables. These distributions are obtained from Monte Carlo studies of single particles passing through the detector.

Clusters which are classified as hadronic must be corrected for the non-compensation of the calorimeter. This is done using a cell weighting approach. However, the approach used for deriving these constants is different to the one used in the ‘GCW’ approach. In this case, dedicated detector simulations store the amount of true energy lost by particles in the different detector regions (absorber layers, active layers, dead material, etc.). This information is referred to as calibration hits. The cell weights are then derived using this information, with the goal of only restoring calorimeter compensation.

Extra corrections for out-of-cluster effects (due to the noise cuts) are applied based on the average distributions from the single particle samples, binned in $\eta$, shower depth and energy. Finally, corrections for energy losses in the dead material are obtained by assigning the energy in dead material regions from calibration hit information to the reconstructed clusters in the simulation. Parameterized first-order functions based on the energy fractions in the calorimeter samplings nearest to each of the sources of energy loss are obtained, with the function parameters derived as a function of cluster $\eta$ and energy.
Early studies have shown that this method is a promising way to improve the detector energy resolution. However, the use of such a method will require in-depth validation of the detector simulation software. Due to the complexity involved with correctly modeling the variables used as inputs to the ‘LC’ calibration, it has therefore not been fully validated for use at ATLAS.

4.4 Jet Cleaning

Spurious or poorly-understood energy deposits in the detector, generated due to various hardware problems, are not removed during the event reconstruction. Therefore, it is possible for fake jets to be formed from such energy deposits.

Studies of early LHC collision data have isolated the presence of various detector faults, leading to either sporadic noise bursts (hot cells), or cells with no signal (dead cells). A set of criteria have been designed to flag jets which are affected by these detector faults [30].

The selection criteria are separated into categories which address two different classes of problems: fake jets, and poorly-understood jets.

4.4.1 Fake Jets

Fake jets are spawned by spurious energy deposits in the calorimeter. While the mechanisms responsible for those is not well-understood, their symptoms are well recognized. Therefore, different selection cuts are targeted at specific detector problems.

The following variables are introduced to address such problems:

- \( n_{90} \): The minimum number of cells required to constitute 90% of a jet’s energy.
- \( f_Q \): The fraction of the energy in the jet associated to cells with a bad \( Q \) factor in the Liquid Argon calorimeter. The \( Q \) factor is a cell-based, \( \chi^2 \)-like measure of the agreement between the observed pulse and the expected pulse shape.
• \( T \): The average time offset between the measured time of the cell signal pulse and the event time for all cells inside a jet. The average is calculated by weighting the time offset of each cell with the square of the cell energy.

• \( f_{HEC} \): The energy fraction of the jet coming from the Hadronic Endcap calorimeter.

• \( f_{EM} \): The energy fraction of the jet coming from the Electromagnetic layers of the calorimeter.

These variables are combined into the following selection cuts used to identify fake jets:

• \( n_{90} < 6, f_{HEC} > 0.8 \): This is used to remove the jets which are caused by sporadic single cell noise bursts observed in the Hadronic Endcap calorimeters.

• \( |f_Q| > 0.8, f_{EM} > 0.95 \): This is used to remove the jets caused by noise bursts correlated across many cells, with bad pulse shapes, in the Liquid Argon barrel calorimeter.

• \( |T| > 50 \text{ ns} \): This is used to remove jets built from out-of-time energy deposits. The 50 ns cut corresponds to a tolerance of two bunch crossings.

Because some of the variables used are not modeled by the detector simulation (for example, only perfect pulse shapes are simulated, and therefore no information on the variability of \( Q \) values can be obtained), it is difficult to obtain a direct estimate of the biases introduced by these selection cuts. However, early studies of the data demonstrated that the acceptance losses triggered by these cuts are negligible. Figure 4.3 shows the effect of applying some of the bad jet cleaning cuts to the dataset selected by the minimum bias trigger. Such a dataset contains the least biased set of events during which an inelastic collision has occurred. The total dataset is shown in blue, with Monte Carlo predictions in red. The gray and black histograms represent the dataset after the \( n_{90} < 6 \),
$f_{HEC} > 0.8$ and $|f_Q| > 0.8$, $f_{EM} > 0.95$ selection cuts are applied, respectively. The spurious $E_T^{\text{miss}}$ tails created by the presence of fake jets are eliminated by applying the quality cuts.

4.4.2 Poorly-Understood Jets

Poorly-understood jets contain real energy deposits located in a poorly instrumented region of the detector (for example, regions where dead cells are present). While these effects will only increase the resolution of the detector, they will introduce tails which are not taken into account by the simulation. It is therefore preferred to ignore events with such jets, at the cost of acceptance, to prevent disagreements with Monte Carlo.

The following variables are introduced to address such situations:

- $f_{COR}$: The fraction of energy from cells which have received the neighboring cell correction in the jet. Cells with recognized permanent problems are masked during the reconstruction. The neighboring cell correction assigns an energy to any masked
cell based on the average energy in all cells by which it is surrounded.

- $f_{TG3}$: The energy fraction of the jet coming from the third layer of the Tile Gap calorimeter.

These variables are combined into the following selection cuts used to identify poorly-understood jets:

- $f_{COR} > 0.5$: This is used to remove the jets which have received a significant correction from the neighboring cell correction, as the effect of the correction, when large, is poorly understood.

- $f_{TG3}$: This is used to remove the jets which have significant energy deposits in the third layer of the Tile Gap calorimeter. This layer is not accessible by the calibration equipment, and therefore, the calibration of the electromagnetic scale in this layer has never been verified.
Chapter 5

Observables and Event Selection

5.1 The Observables

As mentioned in section 2.1, the presence of quark contact interactions would be characterized by an excess of dijet events in the invariant mass regime associated with the contact interaction scale. A direct measure of the dijet cross-section as a function of the invariant mass could be used as a direct probe for such a scenario. However, a mismeasurement of the jet energy scale, which is expected to have a large associated uncertainty for early data, can produce signal-like features in this distribution.

Instead, a set of observables, sensitive to the different angular characteristics of quark contact interactions and QCD processes, are introduced. These observables, typically referred to as angular distributions, have a reduced sensitivity to the jet energy scale, and are therefore more robust for use in early analyses.

5.1.1 \( \chi \)

The most direct way to probe the angular characteristics of dijet final states is to study the scattering angle \( \theta^* \), which represents the angle between the initial state and final state partons in the center-of-mass frame. A direct measurement of \( \frac{d\sigma}{d\cos \theta^*} \) will therefore
provide a simple probe of the angular distributions [6, 7].

However, it has become customary to use the variable $\chi$ in experimental measurements of angular distributions instead [4, 8, 9, 10, 11]. The variable $\chi$ was originally introduced to simplify the expressions for the differential cross-sections of the various QCD $2 \rightarrow 2$ processes [31]. It is defined as a function of the scattering angle $\theta^*$ or, alternatively, as a function of the Mandelstam variables:

$$\chi = \frac{1 + \cos \theta^*}{1 - \cos \theta^*}$$

$$\chi = \frac{\hat{u}}{\hat{t}}$$

where $u$, $t$ and $s$ are the Mandelstam variables, and $\hat{u}$, $\hat{t}$ and $\hat{s}$ are the equivalent center-of-mass variables for the underlying QCD process.

It is important to note that there is nothing fundamental about this choice of observable: it was only chosen because of the many advantages it affords, making it a good choice for experimental analyses. First, it can be demonstrated that the differential cross-section $\frac{d\sigma}{d\chi}$ is expected to vary slowly for Standard Model QCD processes. For the dominant QCD t-channel exchange processes, analogous to Rutherford scattering, the differential cross-section can be expressed as [31]:

$$\frac{d\tilde{\sigma}}{d\tilde{t}} \sim \frac{\alpha_s^2}{\hat{s}^2} F(\chi)$$

where $\frac{d\tilde{\sigma}}{d\tilde{t}}$ is the cross-section of the underlying QCD process, and

$$F(\chi) = \chi^2 + \chi + 1 + \chi^{-1} + \chi^{-2}$$  \hspace{1cm} (5.3)

For massless partons and forward scattering angles characteristic of t-channel processes, the following approximations can be made:

$$\hat{u} + \hat{t} + \hat{s} = 0$$

$$\hat{t} \sim 0$$

$$\hat{u} \sim -\hat{s}$$

$$\chi \sim \frac{\hat{u}}{\hat{t}}$$  \hspace{1cm} (5.4)
Now, expressing the differential cross-section as a function of $\chi$:

\[
\frac{d\hat{\sigma}}{d\chi} \sim \frac{\hat{t}}{\hat{s}} \frac{\alpha_s^2}{\hat{s}^2} F(\chi)
\]
\[
\frac{d\hat{\sigma}}{d\chi} \sim \frac{\hat{t}^2}{\hat{s}^2} \frac{\alpha_s^2}{\hat{s}^2} F(\chi)
\]
\[
\frac{d\hat{\sigma}}{d\chi} \sim \frac{\alpha_s^2}{\hat{s}} \frac{1}{\chi^2} F(\chi)
\]
\[
\frac{d\hat{\sigma}}{d\chi} \sim \frac{\alpha_s^2}{\hat{s}} \left(1 + O\left(\frac{1}{\chi}\right)\right)
\]  

(5.5)

To obtain the final predictions for a hadron collider, the underlying QCD process must be convoluted with the parton distribution functions of the colliding hadrons:

\[
\frac{d\sigma}{d\chi} = \int dx_1 \int dx_2 \ p(f_1, x_1, Q^2) \ p(f_2, x_2, Q^2) \ \frac{d\hat{\sigma}}{d\chi}
\]

(5.6)

where $p$ is the parton probability distribution function for a proton, $f_{1,2}$ is the flavor of the two incoming partons, $x_{1,2}$ are their respective momentum fractions, and $Q^2$ the momentum transfer squared of the event.

Provided that the product of both parton distribution functions remains approximately fixed (by constraining the boost of the event), it can be seen that the differential cross-section is otherwise expected to depend only weakly on $\chi$ (the $O\left(\frac{1}{\chi}\right)$ term indicates that a small rise is expected near the lowest values of $\chi$). Such a behavior has an important consequence for experimental studies: bin migrations due to resolution effects are not expected to introduce drastic changes in the shape of the distribution. This ensures that the observable is more robust against the appearance of fake signals caused by resolution effects.

Another advantage of the $\chi$ observable is that it can be expressed transparently as a function of the final state dijet kinematics:

\[
\chi = e^{\left|y_1 - y_2\right|}
\]

(5.7)

where $y_{1,2}$ are the rapidities of the two leading jets.

Finally, because it is only a function of the center-of-mass scattering angle $\theta^*$, the $\chi$ observable is boost-invariant.
To remove the dependency on the luminosity measurement and the absolute theoretical cross-section, the normalized differential cross-section \( \frac{dN}{d\chi} \) is studied instead, transforming this into a shape-only measurement. Because new physics is expected to produce dijet pairs more isotropically, this would result in an upwards deviation from the expected near flat behavior of \( \frac{dN}{d\chi} \) in the low \( \chi \) regime. The normalized cross-sections are studied in bins of invariant mass. This offers greater sensitivity for signals appearing at a given mass scale, as the QCD contributions from lower kinematic regimes which would dilute the presence of signal are removed.

### 5.1.2 \( R_C \)

The observable \( R_C \), usually referred to as the dijet centrality ratio, is a different probe of the same angular characteristics as \( \chi \). The observable \( R_C \) is defined as the ratio of events where both leading jets are in the inner detector region (defined by \( \eta_1 \)), to the number of events where both jets are in the outer detector region (defined by \( \eta_2 \) and \( \eta_3 \)), as a function of the dijet invariant mass:

\[
R_C = \frac{N(|\eta_{j_1,j_2}| < \eta_1)}{N(\eta_2 < |\eta_{j_1,j_2}| < \eta_3)} \tag{5.8}
\]

where \( \eta_{j_1,j_2} \) is the pseudo-rapidity of both leading jets. Because it is simply a ratio of events, the observable is insensitive to the absolute normalization of the number of events, removing the dependency of the measurement on the integrated luminosity.

The \( R_C \) observable is predicted to be flat as a function of the jet invariant mass for QCD. However, because quark contact interactions are expected to produce more central events, the presence of signal can be observed as a rise in the centrality ratio as the dijet invariant mass approaches the contact interaction scale \( \Lambda \).

While \( \chi \) is a quantity justified by theoretical arguments, \( R_C \) is motivated by experimental considerations. The \( \eta_{1,2,3} \) cuts can be freely chosen to avoid unwanted detector features (detector cracks, different detector technologies, electronic faults, etc.), which
may not be well understood in an early data scenario.

The $R_C$ observable is very robust against fluctuations in the absolute jet energy scale; such variations would only lead to the translation of the $R_C$ distribution as a function of the invariant mass. Because the quantity is expected to be flat for QCD, it is impossible for variations in the absolute jet energy scale to produce a fake signal signature. However, because the $R_C$ observable is a direct ratio of event counts in two detector regions, an incorrect inter-calibration of the detector regions can have significant effects on the distribution.

The $\eta_{1,2,3}$ cuts are chosen to constrain the measurement to a detector region where the inter-calibration is well understood. The barrel region of the calorimeter, up until the transition region, is chosen as the region of focus. In this region, the uncertainty on the inter-calibration (the change in the jet energy scale as a function of the chosen detector region) is known to be of the order of 1% or less. This is an important source of uncertainty for this observable, since event rates in two different detector regions are directly used in conjunction in the ratio. The chosen definition for $R_C$ is then:

$$R_C = \frac{N(|\eta_{j_1,j_2}| < 0.7)}{N(0.7 < |\eta_{j_1,j_2}| < 1.3)} \quad (5.9)$$

The $R_C$ observable is not expected to be as sensitive to contact interactions as $\chi$. This is due to the acceptance losses engendered by the definition of the observable; the requirement that both jets be in the same detector region, along with the smaller pseudo-rapidity range, lowers the acceptance considerably. It is nonetheless an important part of the analysis, as any experimental problem can be traced back to its source much more easily with the $R_C$ observable, due to the detector-driven definition of the observable. Also, due to the simplicity of the observable, it can be studied more finely as a function of the invariant mass. In the case of a discovery, this would allow a better characterization of the observed deviation.
5.2 Kinematic Selection

Specific kinematic selection cuts are implemented for the $\chi$ observable. In early data, the detector acceptance is artificially constrained due to the poor understanding of the jet energy scale in the forward region of the detector. Therefore, only jets within the barrel and endcap regions are considered for this analysis, resulting in a requirement of $|\eta_{\text{jet}}| < 2.8$. Monte Carlo studies have shown that, in the kinematic range studied, using a rapidity requirement of $|y_{\text{jet}}| < 2.45$ ensures that no jets fall outside of the chosen detector acceptance.

A further selection cut is used to define the maximal value of $\chi$ for which $\frac{dN}{d\chi}$ will be measured. This value is directly related to the largest rapidity difference between the two leading jets. While the signals studied here have isotropic angular distributions, and are expected to contribute mostly in the range of small rapidity differences, it is nevertheless interesting to measure the distributions up to large rapidity differences. This allows for a better normalization (because of the greater acceptance) of the spectrum, which in turn leads to better sensitivity to signal. However, as will be detailed in section 5.3.2, a value which is too large will introduce significant biases to the distributions due to trigger effects. The chosen cut is:

$$|y_1 - y_2| < 3.4$$  \hspace{1cm} (5.10)

which leads to a maximal $\chi$ value of $\sim 30$. This value was chosen as a good compromise between acceptance, and the range of masses which can be probed without bias using the trigger with the most statistics.

As shown in figure 5.1, the two selection cuts defined above describe an irregular, diamond-like shape in the rapidity phase space (defined by the variables $y_1$, $y_2$). The resulting acceptance is therefore not uniform as a function of $\chi$. As further illustrated in figure 5.2, the acceptance becomes larger as $\chi$ decreases, consistent with the increase in phase space. Furthermore, the acceptance becomes highly dependent on the kinematic
Figure 5.1: Phase space defined by the various acceptance cuts in \((y_1, y_2)\).

Figure 5.2: Event acceptance as a function of \(\chi\) with (in red) and without (in black) the rapidity boost selection cut.
regime studied. In the regime of larger invariant masses, the contributions to small rapidity differences are restricted to events with small rapidity boosts, which are situated away from the irregularly shaped region. To remove the irregularity in phase space, it is useful to introduce an extra selection cut, which limits the range of rapidity boosts considered. This cut offers the final constraint which leads to a rectangular region in phase space:

\[
y_B < 1.7
\]

\[
y_B = |y_1 + y_2|
\]

(5.11)

where \(y_B\) represents the rapidity boost of the dijet final state. Limiting the range of allowed rapidity boosts also has the advantage of offering a constraint on the different momentum fractions carried by the partons. A tighter constraint on this variable allows more direct probing of the hard-scatter kinematics, by reducing the effects of the convolution of the partonic cross-section with the parton distribution functions.

The binning in \(\chi\) was chosen such that it is linear in rapidity. The detector segmentation leads to a smearing effect on the angular observables. As the detector is segmented using a projective geometry (cells in a detector segment are equally distributed as a function of detector \(\eta\)), the angular resolution is expected to remain of the same order of magnitude as a function of \(\eta\). Therefore, keeping bin boundaries at an equal interval in rapidity (as an approximation for the pseudo-rapidity \(\eta\)) ensures that angular bin migration effects are of similar magnitude across the entire spectrum. The \(\chi\) bin boundaries are therefore positioned at:

\[
\chi_N = e^{(0.3 \times N)}
\]

(5.12)

where \(\chi_N\) is the \(N^{th}\) bin boundary, with the first boundary starting with \(N = 0\). The last bin boundary, \(N = 11\), is artificially extended to 30. The bin intervals of 0.3 were chosen as offering a good compromise between the statistics in a given bin, and a fine granularity in \(\chi\).

For the \(R_C\) observable, the detector regions used to delineate between the central and
non-central region are chosen to be $|\eta_{j_1,j_2}| < 0.7$ and $0.7 < |\eta_{j_1,j_2}| < 1.3$. This region of the detector ends at the transition between the barrel and endcap calorimeters. This ensures that the detector technology remains mostly uniform for the considered region. Preliminary studies showed that the $R_C$ observable is very sensitive to the change in the detector response as a function of $\eta$. Within this range, the relative detector inter-calibration is known to be uniform at the 1% level.

5.3 Event Selection

The data sample considered for this analysis corresponds to an integrated luminosity of approximately $3.1 \, pb^{-1}$ [32]. The events were selected from the $Jet/E_{T}^{miss}$ good run list [30]. This list only contains the runs during which all the detector components needed for this analysis were online and fully functional, and stable beams were present in the accelerator.

The jets used in the analysis were made from topocluster inputs, using the anti-$k_T$ algorithm with a size parameter of 0.6 [33]. Early Monte Carlo studies have shown that such a radius offers good hadronic shower containment for the kinematic range studied, without including a significant contribution from the underlying event.

The leading jet $p_T$ in an event was required to be larger than $60 \, GeV$, and the second leading jet $p_T$ was required to be larger than $30 \, GeV$. These asymmetric cuts were chosen in order to avoid biases when making comparisons with Next-to-Leading-Order (NLO) matrix element calculations. In events where a third jet is radiated, a symmetric cut at the NLO level would result in a more stringent requirement on the first jet. The threshold of $30 \, GeV$ for the second jet was chosen such that the jet reconstruction be fully efficient. Only jets inside the pseudo-rapidity range $|\eta| < 2.8$ were considered for the leading jet selection, all other jets were entirely ignored. This cut was chosen to avoid contributions from the forward region, where the jet energy scale was not fully understood.


5.3.1 Backgrounds and Cleaning Cuts

In this analysis, the events studied have a final state with two jets. The contribution to the set of events with this final state from non-QCD Standard Model processes is negligible since the QCD production rates are known to be orders of magnitude larger. However, non-collision events can constitute a possible source of background to this signature. Such events can come from upstream interactions between the beam particles and gas inside the beampipe (beam-gas) or dead material (beam-halo), cosmic ray muon particles showering in the detector, and fake signals due to detector faults.

To remove such events from the data sample, a well-reconstructed inner detector vertex was required, with at least 5 charged particle tracks matched to it. The vertex was required to be within $|z| < 30$ cm. Such a vertex requirement was shown to be highly efficient in studies with early data.

Also, the spurious jets in non-collision events are expected to possess characteristics which can be used to determine that they were not spawned by a collision (for example, timing information). For this purpose, a set of selection criteria was implemented to deal with such jets. Any jet which is identified as a fake jet (see Section 4.4.1) is removed. Also, to prevent the possibility of mislabeling the leading jets due to large energy fluctuations not modeled by the Monte Carlo simulation, any event with a poorly-understood jet (see Section 4.4.2) with $p_T$ larger than $15 \text{ GeV}$ is rejected.

5.3.2 Trigger Selection

The events used were captured using the inclusive jet trigger chain at level 1, labeled $L1\_JN$, where $N$ corresponds to the electromagnetic scale minimum $E_T$ requirement for a single jet to trigger the event. A study of the inclusive trigger efficiency was conducted as a function of calibrated jet $p_T$, to determine the thresholds where each trigger is assumed to be fully efficient. A trigger is considered fully efficient if the efficiency is higher than
99% and if the acceptance is approximately flat as a function of $p_T$ above this point. Only events with jets above this threshold were used in this analysis; this ensures that the trigger turn-on curves do not have to be modeled, and greatly simplifies the analysis, at a small cost in terms of statistics.

The following triggers were used for the analysis: L1_J15, L1_J30, L1_J90. These were found to be fully efficient for events with at least one jet of $p_T$ larger than 60, 80, 120 $\text{GeV}/c$, respectively. The effective integrated luminosities for these triggers were found to be 0.56, 2.0, and 3.1 pb$^{-1}$.

With the exception of the L1_J90 trigger, which is the lowest unprescaled single jet trigger, all triggers have gradually been prescaled as the instantaneous luminosity increased. This implies that, to make full use of the statistics available, the analysis should be conducted at a kinematic range that is above the point where the lowest unprescaled trigger is fully efficient. This however limits the kinematic range which can be studied. Instead, the binning was chosen such that different kinematic ranges can make use of different triggers to maximize the statistics. This approach is detailed in the following section.

### 5.3.3 Binning in $m_{jj}$

The $p_T$ requirement on the jets introduced by the triggers will cause a bias in the distributions as a function of the invariant mass. The invariant mass of a dijet pair can be expressed as follows:

$$m_{jj} = \sqrt{p_T^1 p_T^2} \sqrt{\chi + \frac{1}{\chi} - 2\cos(\delta \phi)}$$

(5.13)

where $\delta \phi$ is the distance between both jets in the azimuthal plane, and $p_T^{1,2}$ is the transverse momentum of both jets. Assuming that only two balanced jets are present in the event (which is the configuration with the largest mass), the largest invariant mass which
can be obtained is:

\[ m_{jj}^{max} = p_{T}^{\text{trigger}} \sqrt{\chi_{max}^{max} + \frac{1}{\chi_{max}^{max}} - 2\cos(\delta \phi)} \]  \hspace{1cm} (5.14)

Below this mass value, the \( p_{T} \) cut will suppress events with large rapidity differences and lead to a bias in the distributions due to an uneven acceptance as a function of \( \chi \).

Therefore, the choice of \( \chi_{max}^{max} \) determines at which invariant mass an unbiased spectrum can be obtained. For each trigger, \( m_{jj}^{max} \) is found, and can be used to derive the optimal binning to make full use of the available statistics. However, the binning was chosen a priori based on early Monte Carlo studies, and the trigger menu was changed after this study. Therefore, the chosen bins are not optimal, but the correct triggers were still chosen to avoid biases.

For the \( R_{C} \) observable, the same reasoning is used to choose which trigger dataset is used for each invariant mass bin. The expression above simplifies to the following in the simpler case of \( R_{C} \):

\[ m_{jj}^{max} = 2p_{T}^{\text{trigger}} \cosh(\eta_{max}) \]  \hspace{1cm} (5.15)

Bins situated at equal intervals of 50 \( \frac{GeV}{c} \) were chosen for the \( R_{C} \) observable. This decision was made a priori when studying the \( \Lambda \) interpolation (see Section 8.3.1) used in the limit setting procedure. Larger bins would not allow the fitting procedure to converge properly, and smaller bins resulted in very large statistical uncertainties over a large range; the current value was found to be a good compromise. Ideally, a varying bin size matching the detector invariant mass resolution could be used to equalize bin migrations. However, this is not necessary for this analysis since no attempts are made to unfold for detector resolution effects. It is still interesting to note that the bin size is of the same order of magnitude as the detector resolution on the invariant mass for the kinematic range studied.
Chapter 6

Monte Carlo Predictions

The analysis described here consists of measuring the $\frac{dN}{d\chi}$ and $R_C$ distributions from the reconstructed kinematics of the jets in the events. To keep the analysis simple, the observables are defined at ‘detector level’. The jet kinematics are corrected to hadron level, and no attempt is made to unfold the observables to account for fluctuations due to the resolution of the detector. This means that no attempt is made to relate the observables to a parton level, fixed order prediction.

However, a comparison to the standard model expectations is needed in order to form a quantitative statement on the discovery or exclusion of new physics. Such a comparison therefore requires predictions which take into account various effects, both perturbative (parton showers) and non-perturbative (hadronization, detector response, etc.), so that the comparison can be performed on common grounds. This is achieved by using a sophisticated chain of simulations, based on Monte Carlo techniques. This chain will be detailed in the following sections.

6.1 Event Generation

The first step consists of generating the hard scatter process. This is done using the PYTHIA Monte Carlo software [34]. The PYTHIA generator uses Leading Order (LO)
matrix element calculations for the amplitudes for specific physics processes. These calculations are combined with Monte Carlo techniques in order to produce a set of unweighted events with the correct statistical distributions.

The beam particles in the LHC are protons; however, the hard collisions are produced between the constituent partons of the protons. The PYTHIA generator takes this into account by using Parton Distribution Functions (PDFs) to determine the probability to resolve a given constituent parton with momentum fraction $x$ (also known as Bjorken $x$). This also means that the remnants of the scattered protons (beam remnants) are no longer colour neutral, and will interact, with a colour correlation to the remainder of the event. This is taken into account by the underlying event routines of the PYTHIA generator. The radiation which can occur before the constituent partons of the proton interact is also taken into account by the Initial State Radiation (ISR) routine.

To obtain a better prediction of the actual final states observed in a detector, the PYTHIA event generator uses a parton shower approach to create the high parton multiplicity final states observed due to radiative QCD emissions.

Finally, the hadronization of partons is simulated using the Lund string model [34]. This model is empirically motivated; while it does offer a good description of the hadronization procedure (from comparisons of Monte Carlo to data), it is not based on any first principles. A colour string is ‘stretched’ between every pair of partons. As the partons further separate, the energy in the string increases (due to the nature of the strong force) and the string eventually ‘breaks’, producing a new quark anti-quark pair. The procedure stops once the scale is below the hadronization scale and all produced partons can be recombined as colour neutral hadrons.

As a final product, the PYTHIA event generator returns a table of all particles involved in the event, along with their kinematics and relationships.
6.2 Detector Simulation

The GEANT4 software toolkit is used to simulate the passage of particles through the ATLAS detector [35]. This toolkit is a very flexible framework, permitting the full, complex modeling of a particle physics detector, including the correct materials and their physical properties. The software is capable of simulating particle trajectories in magnetic fields, energy losses in materials, particle decays, and various other phenomena relevant to the interactions between particles and a detector. Also, the software is implemented in a way that new physical models of interactions can easily be integrated into the original framework.

The GEANT4 simulation uses a stepping algorithm to simulate particle interactions. For each particle, the list of permitted interactions is probed. From this list of interactions, the mean free path (the average distance before an interaction will occur) is obtained from the combination of all possible interactions. This mean free path is used to determine the step size: the distance the particle will traverse between each iteration of the simulation. At each iteration, the program determines whether an interaction has occurred, using a Monte Carlo approach. The location of the particle is an important ingredient in this calculation, since many of the interactions depend on the material properties of the medium traversed. Therefore, the step size is also dependent on the distance to the nearest material boundary.

By default, GEANT4 offers support for 7 different categories of processes, which can each be implemented by various models. The models used span over 10 orders of magnitude in energy, from 250 eV up to the TeV-scale. The electromagnetic and hadronic categories are perhaps the two most relevant for the description of interactions in the ATLAS detector.

The electromagnetic models are well understood, and have been shown to provide good agreement with experimental data. The models describe a large variety of processes which affect charged particles: bremsstrahlung, ionization, Compton scattering,
A large variety of hadronic models are available in GEANT4 to describe the complex nature of hadronic showers [36]. These models were either derived directly from data, parameterizations, or theory. The default Physics List at ATLAS, which corresponds to a set of chosen models, and the kinematic range over which they are applied, is QGSP_BERT. This list uses the Quark Gluon String model to describe the interactions of hadrons with energies larger than 10 GeV [37]. At lower energies, the Bertini Cascade model is used to describe the interactions between the hadrons and the nucleus in the medium, instead of the default Low Energy Parameterized (LEP) model [38]. The QGSP_BERT list was shown to show the best agreement between shower shapes and detector response in testbeam studies and Monte Carlo simulations [39, 40, 41, 42, 43]. The models in this physics list use a mixture of data tabulation and parameterizations.

6.3 Next-to-Leading Order Predictions

As mentioned in section 6.1, the PYTHIA event generator only calculates matrix elements with two outgoing partons for QCD final states. Final states with more partons are obtained via the parton shower mechanism highlighted in the same section. This process, however, only gives the correct kinematic and angular distributions in the soft/co-linear emission regime. The predictions will not be valid for events where a third, well separated parton carries a significant fraction of the momentum of the collision. Also, the PYTHIA event generator does not include contributions to the $2 \rightarrow 2$ processes from diagrams with loops. These contributions will affect the cross-section contributions from each diagram for QCD processes.

To remedy this problem, Next-to-Leading Order (NLO) event generators include the matrix element contributions for $2 \rightarrow 3$ processes, as well as $2 \rightarrow 2$ diagrams with a single loop. The NLOJet++ event generator is used for this analysis [44, 45]. Ideally,
such a generator should be used as a first step in the full Monte Carlo event generation and simulation chain. However, such NLO generators, due to the techniques used to eliminate divergences in the NLO calculations, produce events with large negative weights. It is therefore difficult to interface such a program to the rest of the event generation chain. Also, it is sometimes still necessary to include parton showers to provide predictions for the high multiplicities of soft/co-linear QCD emissions. This is highly problematic, as both contributions can overlap, and double counting can occur. It is non-trivial to establish a clear separation between the regimes where each approach is preferred.

One possible solution to this problem is known as ‘matching’. In this approach, parton showers are still generated from the final state partons. However, the double counting is avoided by vetoing parton showers which are produced at a higher scale than the scale of the matrix element hard emissions. The vetoing procedure is unique to the different matching schemes which exist. For example, in the MLM scheme, the vetoing is done at the event level, by reconstructing jets from the particle lists, and matching the partons to jets [46]. An event is vetoed if a jet which is not matched to a parton is harder than the softest matrix element parton in the event.

No matching scheme has yet been computationally implemented for QCD jet final states in NLO event generators. Therefore, another approach must be used. The traditional approach consists of studying the effect of the aspects of the simulation which cannot be combined on the full distributions of the observables. Simple, bin-by-bin scale factors are then derived, and applied to the predictions. This assumes that these different aspects of the calculations can be factorized.

To obtain a full prediction, including NLO matrix elements and non-perturbative effects (hadronization, detector simulation), the following approach is traditionally used. The NLOJet++ matrix element distributions are used as the base prediction. Then, a correction factor is used to account for non-perturbative effects. This factor is calculated by comparing the full distributions obtained from a PYTHIA sample which only
includes the matrix elements and parton showers (underlying event, primordial $k_T$, and hadronization turned off), labeled $PYTHIA_{SHOW}$, to a PYTHIA sample (with all parts of the calculation turned on) which has been passed through the full detector simulation, labeled $PYTHIA + FULLSIM$. For a given bin of the distribution, the final result can then be expressed as:

$$\text{FULL} = \frac{NLO_{ME}(PYTHIA + FULLSIM)}{PYTHIA_{SHOW}}$$  \hspace{1cm} (6.1)$$

At the time when this analysis was carried out, no method for calculating NLO predictions for contact interactions had yet been formulated. Because of the inherent complexity in separating the contributions from QCD and contact interactions in the signal samples (due to the presence of interference terms), it was chosen to apply the same NLO correction to both components. This was an arbitrary decision purely aimed at keeping the procedure as simple as possible and ensuring consistency between the QCD predictions in both cases (Standard Model only and Standard Model + contact interactions).

The PYTHIA samples with full detector simulation are therefore used as the base prediction, and are scaled by so called k-factors to account for the better description of the hard parton emissions of the NLO predictions:

$$k = \frac{NLO_{ME}}{PYTHIA_{SHOW}}$$  \hspace{1cm} (6.2)$$

This procedure is mathematically equivalent to the one above, but allows the k-factors calculated from the QCD Monte Carlo samples to be applied to the compositeness samples.

It is interesting to note that recent theoretical efforts have led to the development of an approach to calculate NLO predictions for contact interactions [47]. Such a study showed that an optimal treatment of the NLO corrections to the contact interactions will lower the associated production rates, and therefore lower the observed limits by less than 10%. However, it is important to remember that such a model is only a benchmark, and
that the interaction strength was chosen arbitrarily by historical convention (see section 2.1.1). Most importantly, past Tevatron measurements have applied the NLO correction factors derived from QCD, as was done here. Following a consistent approach therefore allows the comparison of the obtained limits with past experiments on equal grounds. Therefore, the use of such sophisticated NLO calculations for contact interactions is deemed to be unnecessary.

For the $\chi$ observable, the k-factors were found to change the normalized distributions bin contents by up to 6%, with little dependency on the invariant mass. For the $R_C$ observable, the k-factors were not found to significantly differ from unity (compared to the magnitude of the theoretical uncertainties involved in the calculation), and were not applied to the distributions.

### 6.4 QCD Monte Carlo Samples

The Monte Carlo samples used in this analysis have been officially produced using the ATLAS full simulation chain. The PYTHIA generator is used to simulate QCD dijet final states. However, since the production rates for QCD events vary by several orders of magnitude for the kinematic range considered, the number of generated events required to have low statistical uncertainties across the entire kinematic range would be extremely large. Instead, a set of samples covering smaller kinematic ranges are generated with equal numbers of events. When combined, each sample requires to be weighted by its expected cross-section, to recover the original distributions. The variable used to delineate the samples is the $2 \to 2$ hard scatter $p_T$, available in PYTHIA as CKIN(3) and CKIN(4) (the cuts for the lower and upper bounds). These cuts do not translate directly to the final state, since the hard-scatter is generated by PYTHIA before ISR and FSR have been generated, and a safety margin must be included when generating these samples.

A total of 8 samples are generated to describe the full kinematic range. Labeled as
JN, where N is an index from 0 to 8, each sample describes the kinematic regime with the recursively defined lower and upper bounds:

\[
\begin{align*}
p_{\text{min}}^T &= p_T^N \\
p_{\text{max}}^T &= p_T^{N+1} \\
p_T^{N+1} &\sim 2p_T^N
\end{align*}
\] (6.3)

and where \( p_T^0 \) is 7 GeV. Only samples J2 to J6 were used for this analysis (35 to 1120 GeV), since this corresponds to the \( p_T \) range studied in the data.

The process of combining the samples, known as ‘stitching’, is not without difficulties. Due to the variations caused by the ISR and FSR, events near the kinematic boundary of two different samples will receive contributions from both. However, in the sample with the lower kinematic range, these events will be produced very rarely and will therefore have a very large associated statistical uncertainty. However, as the lower kinematic sample has a much higher cross-section, these events will receive large associated weights. This leads to odd fluctuations with large error bars near the kinematic point where two samples are joined, a relic of the stitching process. Various approaches have been suggested to reduce the impact of this effect. One possibility is to implement a per event pre-weighting function, to ensure that events across the full kinematic range are produced at equal rates. This introduces problems of its own, as the weighting function must be carefully optimized a priori to ensure that the correct weight is given to the events based on their importance in a given analysis. Another possibility is to generate samples with no upper bounds. While this addresses the problem, it makes the combination of samples more difficult (as the cross-sections have to be re-calculated based on the chosen cut-off point). However, neither of these schemes have been officially implemented at ATLAS; the default approach was therefore used, leading to the appearance of such odd fluctuations with large associated error bars in some of the distributions.

The generation parameters in PYTHIA are set using the official ATLAS09 tune. This tune was obtained by performing comparisons with data acquired by the Tevatron
experiments [48]. The modified leading-order *MRST2007* parton distribution functions were used as part of the official ATLAS09 tune [49].

Other similar analyses have used other event generators using different techniques to describe the various mechanisms involved. It was shown that the choice of event generator does not have an appreciable effect on the predictions for the kinematic range at which the observables are studied here [50].

The signal samples are also generated using the same ‘stitching’ procedure. The quark compositeness routine in PYTHIA is used. The samples are all generated using destructive interference, as detailed in section 2.1.1. A total of 5 compositeness samples has been generated, covering the range of compositeness scenarios which the analysis was expected to be sensitive to. The 5 samples were generated with a compositeness scale Λ of 500, 750, 1000, 1500, and 3000 GeV.
Chapter 7

Systematic Uncertainties

7.1 Jet Energy Scale Uncertainty

As outlined in section 4.3, a dedicated energy calibration must be applied to jets on top of the electromagnetic scale in order to compensate for the different calorimeter response to hadrons. The uncertainty on this calibration must be studied in detail as this is often the leading source of systematic uncertainty in dijet measurements.

For the first data, the decision made within the ATLAS collaboration was to set a conservative uncertainty on the jet energy scale using combined information from the ATLAS test-beam studies, in-situ studies of the energy response of charged pions, and Monte Carlo simulation studies. The final uncertainties used are derived solely from the latter, with the former two being used as cross-checks and justifications for the decisions taken.

7.1.1 Methodology

Various Monte Carlo samples are generated by varying parameters of the simulation. These variation samples represent conservative estimates of the various systematic effects which could have repercussions on the understanding of the jet energy scale [51].
The response of jets in each sample is measured, using the same matching technique as highlighted in section 4.3.1. The only difference is that no isolation requirement is made, since the uncertainty is to be derived for inclusive jets, which are not always isolated. The response is measured as a function of $\eta^{\text{jet}}$ and $p_T^{\text{jet}}$.

For each variation sample, the response is compared to the response obtained from the nominal sample. The nominal sample represents the best understanding of the underlying physics and of the detector. It is the same sample as used for the analysis, and is described in further details in section 6.4. The difference in the response is taken as the one $\sigma$ uncertainty on the jet energy scale for the given effect studied.

### 7.1.2 Sources of Uncertainty

Figure 7.1 shows a summary of the magnitude of all contributions to the jet energy scale uncertainty. The total uncertainty is shown by the blue band. The specific sources of uncertainty in the jet energy scale will be detailed in the sections below.
Shower Models

The response of the detector is highly dependent on the model used to simulate the passage of hadronic particles in the detector medium, as a change in the shower properties will result in changes in the energy recorded by the active layers of the calorimeter. Samples were generated with two different physics lists in GEANT4: QGSP and FTFP_BERT.

Compared to the nominal choice of QGSP_BERT, in the QGSP model, the Bertini cascades are removed. This leads to showers which are longitudinally shorter. Meanwhile, the FTFP_BERT physics lists substitutes the Quark Gluon String model for nuclear interactions by the Fritjiof model [52].

Studies of the detector response to single pions in the test-beam showed that the differences in the response between these 3 models were larger than the differences between the experimentally observed response and the nominal model [39, 40, 41, 42, 43]. By taking the largest difference in response between the nominal physics list (QGSP_BERT), and either QGSP or FTFP_BERT, a conservative estimate on the uncertainty is obtained.

Detector Description

The density and the effect of the inactive material in front of the calorimeter was not accurately measured, since some of the material was not present during the testbeam. The presence of more or less dead material would affect the energy lost by the hadrons before they reach the active calorimeter layers, and therefore, change the expected response. A sample with a different material budget was generated to estimate the effects of extra inactive material in the detector.

The following changes were made to the default geometry:

- $0.05 \times X_0$ of extra material between the barrel pre-sampler and the calorimeter itself ($|\eta| < 1.45$)
• 0.2 $X_0$ of extra material in the cryostat before the electromagnetic calorimeter ($|\eta| < 1.5$)

• 0.2 $X_0$ of extra material in the cryostat between the electromagnetic barrel and the endcap calorimeter

• 1.5 increase in the density of the material of the barrel-endcap cryostat gap

These estimates come from the combination of testbeam studies and early comparisons between 900 GeV commissioning data and Monte Carlo predictions. Variables such as the relative energy fractions in different calorimeter samplings can be used to study how showers are affected by the extra material, and can therefore be used to obtain an estimate of the extra amount of dead material which may be present.

**Beamspot displacement**

The beamspot, the average center of the collisions in detector coordinates, is not positioned at the center of the detector. While this does not affect the energy measurement of a jet, the pseudo-rapidity will be affected (since jets are reconstructed assuming the detector center as a vertex), and therefore, so will the transverse momentum.

A sample with a wider beam spot displacement of (1.5, 2.5, -9) mm was used to estimate these effects. The generated displacement is a conservative estimate; the average displacement observed during early LHC running was much smaller, at (-0.4, 0.62, -1.3) mm.

**Noise thresholds**

The topoclustering algorithm used to generate the jet inputs is highly dependent on the correct modeling of the detector noise thresholds. Incorrect modeling of the noise could lead to further suppression of real energy deposits, or clustering of detector noise. Two
samples were generated, with the topoclustering algorithm using thresholds 10% lower, and 10% higher, to characterize the effects of changes in the noise modeling.

The 10% figures were chosen based on studies of the stability of the thresholds during dedicated noise runs, and based on the differences between Monte Carlo and real detector noise thresholds.

Electromagnetic Scale

The energy read-out in cells uses a calibration which converts the signal in the calorimeter (ADC counts) to an energy which represents the corresponding average energy deposited by electromagnetic showers (spawned by electrons or photons). This calibration was established during the test-beam.

The jet calibration derived from the Monte Carlo studies assumes a perfect measurement of the electromagnetic scale; any deviations would translate to a direct uncertainty on the jet energy scale.

The uncertainty on the electromagnetic scale for the electromagnetic calorimeter is estimated to be at the 3% level. This estimate is based on studies of the time stability of the calibration of the electronics, and on the differences between the test-beam and full ATLAS setup. For example, the temperature of the liquid Argon in the calorimeter is known to affect the response, and was not precisely measured during the test-beam.

The uncertainty on the electromagnetic scale for the hadronic calorimeter is estimated to be at the 4% level. This estimate was derived by comparing the ratio of response to minimum ionizing muons in data and simulation, for both the test-beam and full ATLAS setup.

These two uncertainties are combined to give a full uncertainty of the jet energy scale based on the average relative sampling of the jet energy in the two separate calorimeters as a function of $p_T^{jet}$.
Fragmentation

The detector response to hadrons is highly energy dependent. Therefore, the fragmentation model used, which, for a given final state parton, determines the nature and energy of the particles in jets which will impact the detector, is expected to affect the response. For example, the detector response to a single pion of a given energy will be much different than the response to two pions of half this energy.

A different set of PYTHIA parameters which affect the fragmentation were obtained using the ‘Professor’ program, and used to generate a variation sample. The Professor program uses a minimization procedure to find the set of parameters which will give the best agreement between a reference data set and the Monte Carlo simulation results [53]. For this particular tune, data from the LEP experiment was used to set the fragmentation parameters.

Underlying Event

A sample which replaces the nominal PYTHIA parameters which govern the underlying event by the set of parameter values referred to as the ‘Perugia’ tune is used to estimate the uncertainty on the underlying event description. The Perugia tune was obtained by tuning the PYTHIA parameters in order to obtain a proper description of the minimum bias event topologies (charged particle densities) observed at the CDF experiment [54]. The effect of this variation is expected to be small, as, in the particle level calibration scheme, the energy due to underlying event is not subtracted from the jet energy, since it still corresponds to particles falling inside the jet area.

Detector Inter-calibration

No estimates of the amplitude of the effects of the mis-modeling of the inactive detector material present in front of the calorimeter were available for the endcap region. Instead, the measure of detector inter-calibration is used to set an upper bound on the effects of
dead material outside the barrel region. In this study, dijet events, where both jets fall in different calorimeter regions, are selected. Since the jets are the only objects present in the event, their transverse momenta are expected to balance. Therefore, the relative scale of different detector regions can be probed [55]. By comparing this relative scale between the data and the Monte Carlo, an uncertainty can be set on the effects of the dead material; if the relative scale is not as predicted, it is assumed to be due to the incorrect modeling of the dead material. Figure 7.2 shows the results of this comparison. Since other effects can lead to such differences, assuming it is solely due to dead material is a conservative overestimate.

It is important to note that, unlike the other uncertainties detailed so far, this uncertainty is assumed to be uncorrelated as a function of $\eta$. For example, the longer showers spanned by the presence of Bertini cascades are expected to increase the response in the entire detector. However, there are no guarantees of such correlations in modeling for the dead material; the material budget could be too low in the central region, and too high in the forward region. For analyses, such as this one, which are very sensitive to the
Chapter 7. Systematic Uncertainties

(a) Jet response as a function of the distance to the nearest jet.

(b) Jet response as a function of $p_T$ for quark and gluon initiated jets.

Figure 7.3: Dependence of the jet response on the flavor and isolation.

angular distribution of the jets, this is an important distinction. The study of detector
inter-calibration can then be used to set a scale on the effects of the decorrelation of the
jet energy scale as a function of detector $\eta$.

**Matrix Element, Parton Showers and Hadronization**

The effect on the jet energy scale due to the choice of Monte Carlo generators used for the
matrix element calculations, parton showers, and hadronization is studied by generating
a dedicated sample. ALPGEN is used for the leading order matrix element calculations
[56]. Processes for $2 \rightarrow N$ final states are generated, for $N$ ranging from 2 to 5. Parton
showering and hadronization are handled using HERWIG [57], and the underlying event
was generated using JIMMY [58].

**Non-closure**

The jet calibrations are derived for isolated jets; this is a requirement to ensure that
the fitting procedure used produces convergent results. However, the signature studied
here is the inclusive jet case. In this case, jets will not necessary be isolated. As the
distance between jets decreases, the response of jets is known to decrease, as shown in
Figure 7.4: Average jet response in $p_T^{\text{jet}}$ and $E^{\text{jet}}$ after the EMJES correction has been applied.

Figure 7.3(a). Closure tests of the calibration demonstrate that the calibration does not bring jets to the correct energy when applied to inclusive jet Monte Carlo samples. This non-closure is taken as an extra, signature-dependent uncertainty.

Also, the calibration is assumed to scale the kinematics directly (assuming all jet constituents would require the same compensation). This may lead to small biases in the measurement of $\eta^{\text{jet}}$, leading to differences between the $p_T^{\text{jet}}$ and $E^{\text{jet}}$ responses. The largest of both is taken as the final uncertainty, to remain conservative. Figure 7.4 shows the response of jets after the calibration has been applied. The deviations from unity are used to set the uncertainty due to the non-closure of the calibration.

Finally, another source of signature dependent non-closure could come from the quark/gluon mixture of jets in the samples. Quark and gluon jets are known to have a different response, as shown in figure 7.3(b), and the calibrations are derived for the quark/gluon fractions expected for dijet events. Since the calibration is derived on the same sample as is used for this analysis, this effect does not come into play here, but would affect a study of, for example, $W +$ jets events.
Extra Interactions

The jet calibration scheme in use for first data makes no attempt to correct for the extra energy included in the jet cones due to multiple interactions. Therefore, the offset in the energy must be taken into account as an extra source of uncertainty.

For this study, the detector is separated into 6400 towers (0.1 by 0.1 in $\eta$ and $\phi$ space). The average energy in each tower is measured as a function of the $\eta_{\text{tower}}$ and of the number of primary vertexes in the event. Since the vertex reconstruction is highly efficient, the number of primary vertexes in an event is considered to be a good approximation of the number of simultaneous interactions occurring.

The difference between the average energy for each tower is calculated, in order to remove the contribution from the underlying event (which is not considered as part of the correction for extra interactions). It is found to scale linearly with every extra primary vertex. Therefore, an energy offset for every added primary vertex can be estimated as a function of $\eta_{\text{tower}}$.

By studying the average number of towers contained in a jet as a function of $\eta_{\text{jet}}$, an estimate on the average energy offset can be obtained as an uncertainty.

7.1.3 Response to Charged Hadrons

The detector response to hadronic showers can be studied in-situ by combining the information from different detector subsystems. The inner detector provides measurements of the transverse momentum of charged particles with high accuracy. This information can be used to study the calorimeter response [59].

Events where a single, isolated, high-quality track is recorded are chosen for this study. Such a selection is aimed at finding events where a single hadron interacts with the detector, and can therefore be used to study the response of single particles. In such events, the track is extrapolated to the calorimeter, in order to find nearby topoclusters.
which should contain the calorimeter signature of the single hadron. Only energy deposits from calorimeter layers where the centroid is close (within $\Delta R$ of 0.2) to the extrapolated tracks are used, to reduce contamination from backgrounds (nearby neutral hadrons which would not leave a track).

By comparing the calorimeter energy ($E$) to the track momentum ($p$), the calorimeter response can be measured. The $\frac{E}{p}$ distribution from data can then be compared to the Monte Carlo distribution to study the accuracy of the detector modeling. For example, the ratio of events where no calorimeter energy deposits are recorded can be used to study the modeling of the dead material, responsible for early showering hadrons which do not reach the calorimeter.

By using the $\frac{E}{p}$ distributions measured from data to fluctuate the reconstructed energy of hadrons in the calorimeter in Monte Carlo dijet samples, an estimate of the shift in the measured energy of jets can be obtained. This shift can be taken directly as the contribution to the jet energy scale uncertainty from the detector modeling and response (dead material, shower models, and EM scale determination). This method is however limited: $\frac{E}{p}$ data was only collected for a limited track $p_T$ range, the study was only carried out in the central region of the calorimeter, and this uncertainty can only be applied to charged hadrons (showers contain other particles, such as neutrons, which are not covered by this analysis). Various extrapolations and approximations must be used to offer an uncertainty which can be applied to jets.

As a consequence, this method was only used as a cross-check. The Monte Carlo based calorimeter-only uncertainty was shown to be consistently higher than the uncertainty derived from single particle studies, as demonstrated in figure 7.5. This is a good indication that the chosen JES uncertainty is conservative.
Figure 7.5: Sum of calorimeter-based contributions to the jet energy scale uncertainty estimated using Monte Carlo methods, and single particle studies.

7.2 PDF Uncertainties

As was described in equation 5.6, the result of a cross-section calculation for a process at a hadron collider is a convolution between the hard scatter differential cross-section, and the parton distribution functions (PDFs). This separation of the short-distance hard scatter and long-distance physics associated with hadrons is known as factorization.

The PDFs are determined using data from various experiments (deep inelastic scattering, jet production, etc.). The different experiments provide complementary data which span a large range of momentum fraction $x$ and momentum transfer $Q^2$ values. A few different collaborations actively provide PDF sets using the up-to-date available data to constantly improve the PDFs. The most widely used PDF sets in the ATLAS collaboration are those provided by the CTEQ and MRST collaborations. All collaborations use similar data and a similar methodology.

The input data is fit using a function of many parameters, to permit interpolation between the many data points. The fitting uses LO or NLO level predictions: the choice of predictions used in the method determines the applicability of the derived PDF sets.
The uncertainty on the determination of the PDFs is directly related to the uncertainty on the fitting procedure. To allow a simple treatment of uncertainties in the PDFs, the fit parameters are diagonalized into eigenvectors to remove any correlations. In the new diagonalized basis, each vector can be related back to a particularity of the PDF: for example, the behavior of a given parton flavor in a specific kinematic range. This procedure allows the different orthogonal eigenvectors to be fluctuated without any need to account for correlations. The PDF sets are provided with ranges of fluctuations along each eigenvector, corresponding to positive and negative one \( \sigma \) variations for each parameter.

The systematic uncertainty on an observable due to the choice of PDFs can then be calculated for both directions according to what is referred to as the Master Equation [60]:

\[
\sigma^+_{PDF} = \sqrt{\sum_i \left[ \max(X_i^+ - X_0, X_0 - X_i^-) \right]^2} \\
\sigma^-_{PDF} = \sqrt{\sum_i \left[ \max(X_0 - X_i^+, X_i^- - X_0) \right]^2}
\]

(7.1)

where \( \sigma^+_/-_{PDF} \) is the upwards/downwards change in the observable due to the PDF uncertainty, \( X_0 \) is the nominal value of any observable, \( X_i^+ \) and \( X_i^- \) are the values of this observable due to one \( \sigma \) shifts in the upwards and downwards direction of eigenvector \( i \), respectively. In other words, this method simply consists in adding in quadrature (since all eigenvectors are uncorrelated) the largest fluctuations in either direction due to the one \( \sigma \) changes in the eigenvectors. For this analysis, the specific observables are the bin contents for the different distributions; therefore, the PDF uncertainties are derived on a bin-per-bin basis.

For this analysis, the CTEQ6.6 NLO PDFs were used in conjunction with NLO-Jet++ to obtain the NLO theoretical prediction. For the LO predictions, the modified MRST2007 PDFs were used, as they are the default for ATLAS. However, the k-factor procedure described in 6.3 ensures that the effects of the choice of LO PDFs are effectively canceled in the ratio. The CTEQ6.6 PDF set includes 22 error eigenvectors, which
were used to calculate the bin-by-bin PDF uncertainties for both observables, using the Master Equation described above [61]. It should be noted that the one $\sigma$ error bands provided by the CTEQ group correspond in fact to 90% confidence intervals, instead of the expected 68%. However, they are traditionally still taken as one $\sigma$ bands.

For certain observables, the differences in distributions attributed to the choice of a particular PDF family was shown in the past to be larger than (or similar to) the differences due to the inherent uncertainties of a given PDF set. Therefore, a traditionally established approach was to use the difference in predictions between different PDF families as an extra uncertainty (this uncertainty is not sufficient by itself as both PDF collaborations use very similar inputs, as mentioned above). However, in the case of these angular distributions, it was shown that the choice of PDF family has a negligible effect on the final predictions. Therefore, it was chosen to derive the final uncertainty based only on the specific PDF set uncertainty eigenvectors.

### 7.3 Scale Uncertainties

The NLO predictions described in section 6.3 have a direct dependency on $Q$, the momentum transfer between partons during the collision. This dependency comes from two different factors. The strong interaction coupling constant, $\alpha_s$, which is used in the matrix element calculation, is dependent on the momentum transfer scale. This dependency is introduced during the renormalization step, and results in the ‘running’ of the constant. The $Q$ scale used for this part of the NLO calculation is referred to as the Renormalization Scale, $\mu_R$. As mentioned in 7.2, the calculation is factorized, separating both the long-distance and short-distance parts of the calculation. The PDFs, which are used to take into account the long-distance physics, are also dependent on the momentum transfer scale. This scale is referred to as the Factorization Scale, $\mu_F$.

For LO, tree level $2\rightarrow2$ processes, the momentum transfer scale can be extracted di-
rectly from the kinematics of the final state, assuming the production mode is known. However, when dealing, more realistically, with $2 \to 3$ processes, calculations including loops, or calculations which include parton showers, the momentum transfer cannot directly be inferred from the final state objects. Therefore, it must be approximated. Since this analysis is concerned by dijet signatures, the momentum transfer is approximated from the final state jets. The following estimator is used:

$$\mu_{R,F} = \frac{p_{T}^{\text{average}}}{} = \frac{p_{T}^{j1} + p_{T}^{j2}}{2}$$  \hspace{1cm} (7.2)$$

where $p_{T}^{j1}, j2$ is the transverse momentum of both leading jets, respectively. In the two or three jet final states considered by the NLO generator, it can be reasoned that the momentum transfer should be well-bracketed by the interval $[0.5 \times p_{T}^{\text{average}}, 2.0 \times p_{T}^{\text{average}}]$.

To account for uncertainties associated with the choice of renormalization and factorization scales, the NLO computations were repeated with the three choices of $Q^2$, $(0.5, 1.0, 2.0) \times p_{T}^{\text{average}}$ varied independently for both scales, for a total of nine possible combinations. The effect of the choice of scales on the observables was therefore computed in a bin-wise fashion, for each of the nine possibilities. This method has traditionally been established as the standard approach for approximating the uncertainties due to the choice of renormalization and factorization scales.

### 7.4 Other Sources

Studies of the choice of Monte Carlo event generator, parton showering and hadronization parameters were carried out in an independent analysis of the similar observables at ATLAS [50]. Such effects were found to be negligible, and were therefore ignored for this study.

Two other sources of uncertainty, the jet energy and angular resolutions, were considered in preliminary studies. While these effects are ultimately taken into account by the detector simulation, no estimates of the systematic uncertainties due to the mis-modeling
of these effects were available when this analysis was carried out. Since then, more recent studies have demonstrated good agreement between the resolutions measured in data and the simulations [62].

Studies emulating the effect of including and excluding the jet energy and angular resolution effects on truth level samples showed no quantifiable differences in the final distributions of the observables. Therefore, assuming that the absolute magnitude of these effects is larger than their relative uncertainty, these sources were found to be negligible, and were ultimately not taken into account.
Chapter 8

Measurement and Limit Setting

8.1 Propagation of Systematic Uncertainties

The effects of the systematic uncertainties on the observables are studied using a pseudo-experiment approach. The dominant systematic uncertainties are listed in section 7. In this approach, the entire analysis chain is repeated on the Monte Carlo predictions for different input conditions, producing different versions of the final distributions. The input conditions defining each pseudo-experiment correspond to different variations of the systematics, generated using the best estimate of the probabilistic distributions of the uncertainties.

8.1.1 Jet Energy Scale

The jet energy scale uncertainties are separated into two components: the absolute scale (average energy scale over the entire detector) and the relative scale (the decorrelation in the scale between the various detector regions studied). The uncertainty probability profiles used as inputs are assumed to be Gaussian and symmetric.

For each pseudo-experiment, a different absolute and relative energy scale is taken from a Gaussian centered at the nominal value (where the calibrations are assumed
to correct the jets to parton level accurately). The width of the Gaussian varies on a per jet basis, as the uncertainty is dependent on the jet transverse momentum and pseudo-rapidity. This particular feature is the main reason for the choice of the pseudo-experiment approach; the jets contributing to each bin in the distributions may have different uncertainties associated to them, and it is therefore impossible to apply the uncertainties naively on a per bin basis.

The uncertainty due to the energy contribution to jets from extra interactions due to pile-up is dependent on the number of collisions which occur in the time span of a single event. Since the Monte Carlo samples do not model the extra interactions, the distribution of the number of primary vertexes (the variable used to estimate the number of simultaneous collisions) is measured in data. The obtained probability distribution function is used to randomly generate a number of extra primary vertexes for every Monte Carlo event in every pseudo-experiment. This randomly generated number is then used as an input to the jet energy scale uncertainty due to pile-up.

It is important to note that the jet energy scale uncertainty can be seen as the experimental uncertainty on the jet energy measurement. It can therefore be reasoned that these uncertainties should be associated with the data distributions. However, since no efforts are made to unfold for the jet energy resolution effects, the comparisons between data and predictions are made at the ‘detector observable’ level. This means that the jet energy scale can instead be attributed to an uncertainty on the modeling of the detector response, and therefore be associated with the theoretical predictions. This approach is preferable since the effects of the jet energy scale on the observables is topology dependent; because the angular distributions of jets are dependent on the theoretical hypothesis, this means that uncertainties, which are dependent on the jet kinematics, will also be. The best example of this topology dependent behavior is in the $R_C$ observable. While the effects of the absolute jet energy scale are nearly negligible on a bin-by-bin basis for QCD because variations in the absolute energy result in the translation of a flat
line, they become non-negligible in the case of signal, where the prediction is no longer flat.

### 8.1.2 PDF Uncertainties

The uncertainties associated with the PDFs have been calculated for each observable on a bin-per-bin basis, as highlighted in section 7.2. For each pseudo-experiment, a number $\delta_{PDF}$ is generated from a Gaussian centered at 0 with a $\sigma$ of 1. This number is then multiplied by $\sigma_{PDF}^+$ or $\sigma_{PDF}^-$, the bin-by-bin asymmetric uncertainties of the PDF, depending on the sign of $\delta_{PDF}$. The final value is applied as a variation to every bin in each distribution. This procedure is followed to ensure that the asymmetry characteristic of the derived PDF uncertainties is properly taken into account. The PDF uncertainties are considered as uncertainties on the Monte Carlo predictions.

### 8.1.3 Scale Uncertainties

Because no knowledge of the expected distributions of the uncertainties due to the choice of factorization and renormalization scale for the NLO calculations is available, a flat prior is used. This means that the 9 different NLO distributions, obtained as described in section 7.3, are used as seeds for the pseudo-experiments, and are therefore considered equally likely. The scale uncertainties are considered as uncertainties on the Monte Carlo predictions.

### 8.1.4 Convolution of Uncertainties

A set of 1000 pseudo-experiments are used to model fluctuations of the jet energy scale and PDF, with each pseudo-experiment randomly generating a different scenario for both uncertainties. Each of these pseudo-experiments is used with a different seed for the 9 different choices of scales, as each choice is considered equally likely. This results in
a total of 9000 pseudo-experiments. This number was chosen as it is a good trade-off between processing speed and stability of the final results. Studies have shown that further increasing the number of pseudo-experiments does not change the final results appreciably. However, since each pseudo-experiment requires a full iteration of the event processing chain, it is preferable to keep the final number as low as possible.

To produce the uncertainty bands used for comparisons with Monte Carlo in the final measurement, the 9000 final distributions resulting from the pseudo-experiments are used. For each bin, the central value is defined by the nominal expected values for the Jet Energy Scale (jets at hadronic scale with no detector decorrelation), choice of PDF (central member) and choice of scale for the NLO calculations ($p_T^{\text{average}}$ for both the renormalization and factorization scales). For each bin, a likelihood distribution of the bin contents is obtained from the 9000 pseudo-experiments. The likelihood is integrated in a single direction, starting from the central value, until the 68% interval is found; this interval corresponds to the estimate of the one $\sigma$ deviation in the given direction. During the integration, the regions above and below the central value are treated separately, correctly yielding asymmetric uncertainties.

Figure 8.1 shows a summary of all systematics as a function of $m_{jj}$ for the $R_C$ observable, and for the highest invariant mass bin ($m_{jj} > 1200$ GeV$^2$) in $\chi$.

### 8.2 The Measurement

#### 8.2.1 Evaluation of Statistical Uncertainties

Since the $\chi$ observable consists of a direct measurement of $\frac{dN}{d\chi}$, the uncertainties can be obtained directly from the Poisson distribution function.

The procedure for $R_C$ is not as simple. Past experiments have used the assumption that the likelihood for an observable using the ratio of two event counts ($N_{\text{inner}}$ and $N_{\text{outer}}$}
in this case) can be approximated as a Gaussian, with the associated error propagation:

\[
\frac{\Delta R_C}{R_C} = \sqrt{\left(\frac{\sqrt{N_{inner}}}{N_{inner}}\right)^2 + \left(\frac{\sqrt{N_{outer}}}{N_{outer}}\right)^2}
\]  \tag{8.1}

However, for the number of events present in the bins where the sensitivity to new physics is the highest (in the order of 1 to 2 events in the highest \(m_{jj}\) bins), it was found that this approximation is ill-justified. The Poisson distribution functions are highly asymmetric for small numbers of events, which should yield equally asymmetric uncertainty profiles for the \(R_C\) observable. Since no closed form expression can be derived to express the expected uncertainties, a likelihood profile approach is used instead. For each bin, the phase space of possible \(R_C^{\text{theory}}\) values is swept, with the likelihood that \(R_C^{\text{data}}\) is a statistical fluctuation being evaluated. The profile is integrated around the central value in both directions until the 68% interval is constructed to obtain the statistical uncertainties. This yields highly asymmetric uncertainties in the bins with low statistics, as is expected.
8.2.2 Results

Figures 8.2(a) and 8.2(b) show the comparison between QCD Monte Carlo predictions and data for the $R_C$ and $\chi$ observables, respectively [63]. The data points include statistical uncertainties only, evaluated using the techniques described in section 8.2.1. The MC predictions include the systematic uncertainties from theoretical sources (PDF uncertainties, Scale uncertainties, Monte Carlo statistics) in yellow, and systematic uncertainties from experimental sources (JES) in orange, as described in section 8.1. For the $\chi$ observable (figure 8.2(b)), the distributions for the four invariant mass bins studied are shown simultaneously on the same plot. Since they are all normalized to unit area, a fixed offset was added to each distribution in order to improve visibility. For the $R_C$ observable (figure 8.2(a)), the last bin in $m_{jj}$ is inclusive, and contains events up to 2.5 TeV.

In both figures, the expected distributions for a sample compositeness hypothesis were shown in dotted lines, to demonstrate the sensitivity of the observable to the presence of signal. It is interesting to note that, for the $R_C$ observable shown in figure 8.2(a), the data favors a compositeness scale $\Lambda = 2.9$ TeV, as evidenced by the apparent rise in $R_C$ at high invariant mass. However, further studies show that the data is still compatible with the QCD predictions within a 68% shortest interval.

8.2.3 Agreement with Monte Carlo

Before limits on the existence of new physics can be set, the agreement between the measured data and the Standard Model theoretical predictions must be quantified. A $\chi^2$-like figure of merit is used, using the following definition:

$$\chi^2 = \sum_{\text{all bins}} \left( \frac{\text{data} - \text{theory}}{\Delta \text{data}} \right)^2$$

where $\Delta \text{data}$ is the statistical uncertainty on the measured data for a given bin.

This figure of merit differs from the traditional use of the $\chi^2$ test in its interpretation, since the systematic uncertainties cannot be assumed to be uncorrelated across bins. Sets
of pseudo-experiments are generated, varying the systematic uncertainties following the method highlighted in section 8.1.4 while keeping track of all correlations. A likelihood profile of the obtained $\chi^2$ values is generated for each observable. These profiles are then used to extract a p-value for each observable, which quantifies the probability that fluctuations due to the systematic uncertainties are larger than the ones seen in the data.

For the $\chi$ observable, the p-values are found to be 0.19, 0.11, 0.27 and 0.54, for the lowest to highest invariant mass bins. For the $R_C$ observable, the p-value is found to be 0.85. From these p-values, good agreement between the Standard Model and the measured data can be inferred.
8.3 Setting Limits

8.3.1 Interpolation in $\Lambda$

A set of Monte Carlo full simulation samples was generated for different compositeness scales $\Lambda$. Due to the limited amount of computing resources available for the production of these samples, only a small set of compositeness scales could be generated. Five samples were generated with $\Lambda$ values of 500, 750, 1000, 1500 and 3000 GeV, sampling the range of possible hypotheses which can be studied with the data sample used.

In order to set more precise limits on the exclusion of contact interactions, a method for interpolation in $\Lambda$ was devised for each distribution under study. This is particularly important in the Bayesian limit setting case, since the limits are obtained by integrating the posterior probability function over the range of priors, which are themselves dependent on $\Lambda$. If no interpolation was used, the small number of compositeness Monte Carlo samples available would yield a very coarsely defined posterior curve.

In the case of the $R_C$ observable, the final distribution for each $\Lambda$ value is fit using
a modified Fermi function, which was chosen since it modeled the characteristics of the compositeness sample distributions well; it possesses two flat regions (plateaus) separated by a rising slope (turn-on). These fits are shown in Fig. 8.3(a). The modified Fermi function is defined as:

\[ f(m_{jj}) = \frac{p_1}{\exp\left(\frac{p_3 - m_{jj}}{p_4}\right) + 1} + p_2 \]  

(8.3)

where \( p_1 \) is the difference in amplitude between the two plateaus, \( p_2 \) is the value of the first plateau, \( p_3 \) is the turn-on point and \( p_4 \) characterizes the slope of the turn-on.

By fitting the obtained parameters \( p_i \) as a function of \( \Lambda \), a functional which can be used to describe the \( R_C \) distributions as a function of \( m_{jj} \) and \( \Lambda \) is obtained. Studies of these fits show that the parameters associated with the plateaus (\( p_1 \) and \( p_2 \)) do not vary appreciably and are kept constant. The remaining parameters \( p_3 \) and \( p_4 \) are fit with polynomials of the first and second degree with good agreement found in both cases. The choice of the polynomial used was found not to affect the final limits. This functional is used as the signal hypothesis during the remainder of the limit setting procedure. Note that the \( \Lambda = 500 \) GeV distribution does not follow the same behavior as the others; this is due to the fact that the matrix element compositeness turn-on occurs before the turn-on due to the PDFs for quarks. Therefore, the \( \Lambda = 500 \) GeV point is not used in the interpolation and limit setting procedure. This causes no problems as the limits are set in a much higher kinematic range, where the Fermi function remains a good approximation for the distributions.

No function of \( \Lambda \) which could be used as a parameterization of the full \( \chi \) distributions was found. Instead, a separate parameterization was used for every single \( \chi \) bin, modeling the change in the normalized differential cross-section as a function of \( \Lambda \) on a per-bin basis. The modified Fermi function described above was used to perform the 11 fits for the \( \chi \) bins, as it was empirically demonstrated to serve as a good model for the differential cross-section. Figure 8.3(b) shows a sample fit for the first bin in \( \chi \).
8.3.2 Frequentist Limit Setting using $\chi$

A frequentist method was chosen as the main method for limit setting. Prior independence and simplicity are the main reasons for this decision. The ratio of events contained in the first four bins to the total number of events is used as the observable, defined as $F_{\chi}$. This collapses the entire $\chi$ spectrum into a single number, which greatly simplifies the limit setting procedure. A priori studies of $F_{\chi}$ in Monte Carlo simulations were used to optimize the bins chosen such that $F_{\chi}$ showed a sensitivity comparable to methods using the shape information from the full spectrum.

To set limits using the frequentist method, a large set of pseudo-experiments are generated for each signal hypothesis. First, a value of $F_{\chi}$ is taken from one of the 9000 systematic variations. Then, the number of events in the first 4 bins (the numerator of $F_{\chi}$) is randomly fluctuated using Poisson statistics, under the constraint that $N_{\text{total}}$ (the denominator of $F_{\chi}$) is the same as what was observed in data. This yields the final value of $F_{\chi}$ for a given pseudo-experiment. This ensures that, for each pseudo-experiment, the value of $F_{\chi}$ takes into account possible fluctuations due to both systematic and statistical

Figure 8.4: Neyman construction used in frequentist limit setting for the $\chi$ observable.
For each signal hypothesis, a binned likelihood distribution for the variable $F_\chi$ is generated. Each likelihood is then integrated in one direction, towards the $F_\chi$ value for the Standard Model expectation, until 95% of the area is included. This point is the observed value of $F_\chi$ required to set 95% C.L. limits on this signal hypothesis. The expected limit is therefore obtained for the model where the 95% one-sided point corresponds to the Standard Model prediction, and the observed limit where it corresponds to the measured value of $F_\chi$. A graphical construction which allows the visualization of this procedure is referred to as a Neyman plot. The model parameter is displayed on the x-axis and the value of $F_\chi$ on the y-axis. The intersection between the horizontal line drawn for the observed value of $F_\chi$ and the curve which shows the 95% one-sided value of $F_\chi$ as a function of the model parameter shows the extracted limit. This construction is shown in Fig. 8.4.

### 8.3.3 Bayesian Limit Setting using $\chi$

As an alternative to using the frequentist method, limits are also set using the Bayesian approach. While more complicated, this approach allows the information from the entire $\chi$ spectrum to be used. In this approach, Bayes’ theorem is used to find the posterior probability $P(t|d)$; the probability that the data ($d$) observed represents a given theory ($t$):

$$P(t|d) = \frac{P(d|t)P(t)}{P(d)}$$

(8.4)

where $P(d|t)$ is the probability that the data was generated by a given theory, and $P(t)$ is known as the prior probability distribution function. The prior is a subjective choice of the expected probability that a given theory exists. Since $P(d)$ is independent of the theory, it is traditionally omitted from the expression altogether. In the limit setting procedure, $P(d)$ simply contributes to a normalization factor which will not affect the
limits obtained.

The limits are calculated by integrating the posterior over the space of priors corresponding to all theories until the point where the 95% C.L. is obtained in the likelihood integral:

$$\int_{\xi_{\text{min}}}^{\xi_{\text{95\%}}} P(t|d) \, d\xi = 0.95 \int_{\xi_{\text{min}}}^{\xi_{\text{max}}} P(t|d) \, d\xi$$  \hspace{1cm} (8.5)

where $\xi$ is a parameter of the theory, $\xi_{\text{min}}$ is the value at which the hypothesis corresponds to the standard model, $\xi_{\text{max}}$ is the upper limit of the model parameter, and $\xi_{95\%}$ is the 95% C.L. limit set on this parameter.

In the case of quark contact interactions, the model can be entirely described by the compositeness scale parameter $\Lambda$. In order to set limits, a prior must be chosen. The conventional approach for such searches is to assume that the probability for the existence of a model is proportional to its cross-section. In the case of contact interactions, this corresponds to a flat prior in $\frac{1}{\Lambda^4}$. As a cross-check, a prior proportional to the interaction strength (flat in $\frac{1}{\Lambda^2}$) is also studied. In the case of contact interactions, the expression to be solved becomes:

$$\int_{0}^{\Lambda^{-4}_{95\%}} P(t|d) \, d\Lambda^{-4} = 0.95 \int_{0}^{\infty} P(t|d) \, d\Lambda^{-4}$$  \hspace{1cm} (8.6)

For the $\chi$ observable, the $P(d|t)$ term corresponds to the product of the probability likelihood for each bin of the observable:

$$P(d|t_i) = \prod_{\text{all bins}} \mathcal{L}_i$$  \hspace{1cm} (8.7)

Each theoretical prediction is dependent on the choice of pseudo-experiment, described in section 8.1.4, represented by the index $i$. These pseudo-experiments account for the changes in the predictions due to the effects of the systematic uncertainties. This dependency is marginalized by summing the probability distribution function over all pseudo-experiments:

$$P(d|t) = \sum_i \prod_{\text{all bins}} \mathcal{L}_i$$  \hspace{1cm} (8.8)
The likelihood for a given bin is calculated as follows, using the fact that the probability to observe \( N^{data} \) events in a bin is governed by Poisson probabilities:

\[
\mathcal{L} = \mathcal{P} \left( N^{data} \mid N^{theory} \right) \left( \sum_{\text{all bins}} N^{data} \div \sum_{\text{all bins}} N^{theory} \right)
\]

where \( \mathcal{P} \) is the Poisson probability distribution function. By normalizing the theoretical predictions to the total number of events observed in data, the measurement becomes solely dependent on the shape of the distributions: the predicted absolute production rates and luminosity information are no longer necessary.

The expected limits can be obtained by applying the same limit setting procedure to the expected Standard Model distributions instead of the observed spectrum, while keeping the original constraint on the total number of events observed.

### 8.3.4 Bayesian Limit Setting using \( R_C \)

A Bayesian approach is used in order to set limits with the \( R_C \) observable, as previously described. The procedure used is identical, the only difference is in the expression used to calculate the likelihood, which is dependent on the observable. For a given theory hypothesis defined by the choice of \( \Lambda \), the probability likelihood is calculated for each bin using the expression:

\[
\mathcal{L} = \mathcal{P} \left( N^{data}_{\text{inner}} \mid N^{theory}_{\text{inner}} \right) \mathcal{P} \left( N^{data}_{\text{outer}} \mid N^{theory}_{\text{outer}} \right)
\]

where \( \mathcal{L} \) is the likelihood for a given bin, \( \mathcal{P} \) is the Poisson probability distribution function, \( N_{\text{inner}} \) and \( N_{\text{outer}} \) are the numbers of events in the inner and outer regions, respectively, and they add up to \( N_{\text{total}} \). Because \( R_C \) is a shape measurement, no rate information from the Monte Carlo predictions is used in the expression above. The total measured number of events in each bin, \( N^{data}_{\text{total}} \) is used as a normalization constraint to calculate the predicted number of events for the theory, yielding a different expression which is only dependent on the theoretical prediction of \( R_C \):


\[ \mathcal{L} = \mathcal{P} \left( \frac{N_{\text{data}}^{\text{inner}}}{N_{\text{data}}^{\text{total}}} \left| \frac{R_C^{\text{theory}}}{1 + R_C^{\text{theory}}} \right. \right) \mathcal{P} \left( \frac{N_{\text{data}}^{\text{outer}}}{N_{\text{data}}^{\text{total}}} \left| \frac{N_{\text{data}}^{\text{total}}}{1 + R_C^{\text{theory}}} \right. \right) \]  

\[(8.11)\]

### 8.3.5 Results

Using the methods described in the above sections, limits are set on the existence of quark compositeness at a scale \( \Lambda \). Using a frequentist approach with the \( \chi \) observable in the highest invariant mass bin, 95\% C.L. limits are set at \( \Lambda = 3.4 \text{ TeV} \), as shown in figure 8.4. The expected limits found at the point where the Monte Carlo predictions cross the 95 \% C.L. in the figure, are at \( \Lambda = 3.5 \text{ TeV} \). Using the same observable, a confirming analysis is performed using the Bayesian approach. As described in section 8.3.3, two different choices of priors were used. The 95 \% C.L. limits are found to be 3.2 and 3.3 TeV, for priors flat in \( \frac{1}{\Lambda^4} \) and \( \frac{1}{\Lambda^2} \), respectively. A sample posterior probability distribution is shown in figure 8.5(b), with the red line indicating the 95\% C.L. limit on \( \Lambda \), with a prior of \( \frac{1}{\Lambda^4} \).

For the \( R_C \) observable, a Bayesian approach was chosen, with the same 2 choices of
prior as for the $\chi$ observable. A sample posterior probability distribution is shown in figure 8.5(a), with the red line indicating the 95% C.L. limit on $\Lambda$, with a prior of $\frac{1}{\Lambda^4}$. The 95% C.L. limits are 2.0 and 2.4 TeV, respectively for priors flat in $\frac{1}{\Lambda^4}$ and $\frac{1}{\Lambda^2}$. The corresponding expected limits are of 2.4 and 2.5 TeV.

The limits on contact interactions set by the frequentist $F_\chi$ analysis can be converted to a corresponding length scale at which the structure of the quark has been probed. The 3.4 TeV limit corresponds to a length scale of $6 \times 10^{-5}$ fm.
Chapter 9

Conclusion

9.1 Summary

The results of the first study of the dijet angular distributions at ATLAS were presented. The angular distributions were probed using two different observables established in past experiments: the $\chi$ distributions and the dijet centrality ratio $R_C$. The $\chi$ approach has proven to be more sensitive to the presence of new physics, while the $R_C$ observable has been shown to be a simpler, more robust quantity, very well suited to early data analyses.

Good agreement was observed between the data and Monte Carlo QCD predictions. New limits have been set on the existence of quark compositeness, surpassing the limits set by past experiments. The limits using the frequentist approach were set at $\Lambda > 3.4$ TeV, while the latest limits from the Tevatron were set at $\Lambda > 3.1$ TeV. The current limits are equivalent to a length scale of $6 \times 10^{-5}$ fm. Limits at 4 TeV were set by the CMS collaboration with a similar dataset [64]; however, the expected limits obtained by the CMS and ATLAS experiments are comparable.
9.2 Future Prospects

As no evidence of quark compositeness was found in this analysis, it is crucial to repeat such a measurement as more integrated luminosity is acquired. Studies of the limits as a function of the integrated luminosity have shown that the analysis is currently in a regime where the results are statistically driven: the statistical reach in invariant mass is the main factor determining how high limits on the compositeness scale can be set. By diminishing the magnitude of the systematic uncertainties, particularly the absolute jet energy scale, the limits can be further improved at the few percent level. Such a gain is minimal compared to an increase in integrated luminosity, but will nonetheless improve the sensitivity of the analysis for fixed size datasets.

The experience and knowledge acquired during the study of the two angular observables in this analysis has since led to the development of a new observable, $F_\chi$. By collapsing the entire $\chi$ distribution into a single number (as was done in the frequentist limit setting, described in section 8.3.2), it becomes possible to greatly increase the granularity in invariant mass. This amounts to combining the large acceptance and boost invariance of the $\chi$ observable and the fine mass binning of the centrality ratio $R_C$. Such an observable proves to be more sensitive to new physics, since the coarse invariant mass binning used for the $\chi$ observable resulted in signal dilution.

Finally, such an analysis can also be used to set limits on other exotic physics signals. The angular distributions are sensitive to any signal which possesses different angular characteristics than t-channel QCD processes. For example, excited quark models produce resonant dijet pairs at a scale $M$ via s-channel production. Such a venue is currently being explored as an extension to this analysis.
Bibliography


