Master thesis
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Measurement of Z bosons
in p+Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV with the ATLAS detector

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Abstract

The aim of this thesis is to study the geometry of heavy ion collisions and nuclear effects.

This thesis describes the measurement of Z bosons produced in p+Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV using the ATLAS detector at the LHC. Z bosons - produced by the Drell-Yan process - are measured in the muonic decay channel. From an integrated luminosity of $25.5 \text{ nb}^{-1}$, 1941 Z bosons are observed with a background contamination, estimated from the number of same-sign dimuons, of 0.3%.

The measurement is performed within different centrality groups. From the distribution of the yield within each centrality group, Z production is found not to scale with the number of binary collisions using the default Glauber Model. Instead it is shown, that a correct description of the centrality in p+Pb collisions might require the Glauber-Gribov Model, which includes fluctuations in the nucleon-nucleon cross-section.

Apart from the centrality determination, the differential $Z \rightarrow \mu\mu$ cross-section is measured. The differential cross-section is found to be:

$$\sigma = 182^{+15.1}_{-13.4} \text{ nb}$$

No earlier measurements of Z production in p+A collisions exist, and direct comparison is therefore not performed. The comparison is instead done with the cross-section obtained in p+p collisions assuming binary scaling. It gives an enhancement of Z production in p+Pb collisions of about 50% compared to Z production in p+p collisions, which cannot be explained by well-know effects.
# Contents

1 Introduction

2 The Standard Model
   2.1 Particles and interactions of the Standard Model
       2.1.1 Fermions
       2.1.2 The gauge bosons
       2.1.3 Feynman diagrams
       2.1.4 Transition probabilities
   2.2 Quantum Field Theory
       2.2.1 The Lagrange formulation
       2.2.2 Gauge invariance
       2.2.3 The Higgs mechanism
   2.3 Electroweak Unification
   2.4 Quantum Chromodynamics
       2.4.1 Divergencies and Renormalisation
       2.4.2 Running couplings and asymptotic freedom

3 Probing the hadronic structure and examining the $Z$ boson
   3.1 The structure of the proton
   3.2 Drell-Yan production
   3.3 The $Z$ boson
   3.4 Kinematics
       3.4.1 The ATLAS coordinate system
       3.4.2 Transverse momentum
       3.4.3 Rapidity

4 Heavy Ion Collisions
   4.1 Multiplicity
   4.2 Geometry of a heavy ion collision
   4.3 Glauber Model and centrality
       4.3.1 Relating Glauber parameters to data
       4.3.2 Relating $N_{coll}$ and $N_{part}$
   4.4 Nuclear phases
       4.4.1 Quark Gluon Plasma
       4.4.2 Color Glass Condensate
   4.5 Nuclear PDFs
   4.6 Binary scaling?
5 The LHC and the ATLAS detector  
5.1 The ATLAS detector  
5.1.1 The inner detector  
5.1.2 The calorimeters  
5.1.3 The muon spectrometer  
5.1.4 The trigger system  
5.1.5 Minimum Bias and muon triggers  
5.1.6 Muon reconstructions algorithms  

6 Measurement of the Z boson  
6.1 Data samples  
6.1.1 Monte Carlo  
6.2 Event selection  
6.2.1 Minimum bias triggers  
6.2.2 Pileup  
6.2.3 Diffractive events  
6.2.4 Criteria for the selection of Z bosons  
6.3 Background estimation  
6.3.1 Background sources  
6.3.2 Background estimation from same-sign dimuons  
6.4 Optimisation of the offline cuts  
6.5 Background corrected plots  
6.6 Centrality dependence  
6.6.1 Centrality in the data  
6.6.2 Correction of FCal E_T  
6.6.3 Centrality plots  
6.7 Efficiencies and acceptance  
6.7.1 Trigger efficiencies  
6.7.2 Reconstruction efficiencies  
6.7.3 Acceptance  
6.8 Cross Section  

7 Conclusion  

Bibliography
Chapter 1

Introduction

The field of high energy physics is dedicated to the study of elementary particles and the interactions between them, in order to address the question; "What is matter made of?" at the most fundamental level.

This question is not new and has probably created curiosity amongst people at all times. The idea that matter consist of small elements has been known since the ancient Greece, where Democritus formulated what is though to be the first atomic theory [1]. The word "atom" - meaning indivisible - was later adopted by the atom described by Niels Bohr [2]. In 1897, when the English physicist J.J. Thomson discovered the electron and proposed a model for the structure of the atom, it was found to be composed [3]. Some will say, that by this discovery the area of particle physics was born [4].

Since the end of the 20th century, the knowledge of particle physics has been improved tremendously. Experimentally, the field has evolved from low-energy scattering experiments to collision experiments at very large energies. As a result of this improvement a wide range of elementary particles and their interactions has been discovered. This collection of particles and interactions is described by the Standard Model of particle physics. Despite the comprehensive knowledge already achieved, a lot of questions still need to be answered. What causes the gravitational force? Why does matter dominate antimatter? How did the Universe look like at the Bang Bang? etc.

By the study of high energy collisions, some of these questions could be answered. Collisions between high energetic particles took its beginning during the 60s and 70s. In the 80s two important colliders were ready for operation; The Tevatron at Fermilab and LEP at CERN. LEP is now replaced by the LHC operating at energies of 7 TeV. Proton-proton or positron-electron collisions have been and are still the dominating types of collisions. However, collisions between heavier nuclei have also been studied.

In 2010, a new program for heavy ion physics was started at the LHC. This program made it possible to study Quark Gluon Plasma, which is a hot and dense state of matter emerging from colliding heavy nuclei. This state is also believed to
CHAPTER 1. INTRODUCTION

describe the earliest state of the Universe. Heavy ion collisions gives therefore the possibility for an experimental study of this area. However, heavy ion collisions are not only characterised by the Quark Gluon Plasma, but also by additional nuclear effects caused by a phase of matter known as Color Glass Condensate. To study these effects and to decouple nuclear effects from the Quark Gluon Plasma, collisions are made both between proton and lead and between lead and lead.

This thesis describes the measurement of $Z$ bosons produced in $p+Pb$ collisions at $\sqrt{s_{NN}}=5.02$ TeV using the ATLAS detector at the LHC. The study includes a description of the collision geometry and what is characterised as central collisions compared to peripheral collisions. In addition, different nuclear effects are presented and the impact of the centrality of the collision is discussed. The measurement of $Z$ bosons, which is done via the muonic decay channel, is performed for different centralities and the distribution of the yield is discussed, taken geometric and nuclear effects into account. In addition the differential $Z \rightarrow \mu\mu$ cross-section is determined.

The first three parts of this thesis describe the theoretical framework and the expectations. First the particles and interaction of the Standard Model are introduced. Followed by a description of the proton structure, the Drell-Yan production of the $Z$ boson and a definition of a number of useful kinematic variables. The last theoretical part addresses heavy ion collisions. It describes the geometry of the collision including a description of the Glauber Model, the state of matter, which is formed, and the nuclear effects that appear.

The theoretical part is followed by a description of the ATLAS detector, which is used to measure the muons decaying from the $Z$ boson.

At the end, the analysis is presented. Apart from the measurement of the $Z$ bosons at different centralities, it also includes data driven detector efficiency determination. This is a necessary part to estimate the total number of produced $Z$ bosons and thereby be able to measure the differential $Z \rightarrow \mu\mu$ cross-section.
Chapter 2

The Standard Model

Everything we observe around us is formed by elementary particles - as the name suggest, the most fundamental constituents of the universe to the extend of our current knowledge. Elementary particles and their interactions are described by the Standard Model of particle physics [4]. It was formulated during the 1960s and 1970s on the basis of two families of particles - quarks and leptons - by incorporating Electrodynamics, the Glashow-Weinberg-Salam Theory of the electroweak interaction and Quantum Chromodynamics. The Standard Model describes all known elementary particle interactions except from gravity.

2.1 Particles and interactions of the Standard Model

The Standard Model of particles describes fermions and gauge bosons, known as matter particles and force carriers, respectively.

Fermions include two families - quarks and leptons - while gauge bosons include photons, gluons, Goldstone bosons and the Higgs boson. The Standard Model describes three different kinds of interactions: the strong, the electromagnetic and the weak interaction. The strong interactions are mediated by gluons, the electromagnetic interactions by photons and the weak interactions by Goldstone bosons, which count charged $W^\pm$ bosons and neutral $Z$ bosons. The Higgs boson is a quite special particle, which is mediated between all particles - including itself. All these particles and their properties are illustrated in figure 2.1.

2.1.1 Fermions

All fermions have spin 1/2 and obey Fermi-Dirac statistics. Both quarks and leptons exist in three generations. The two lightest particles make up the first generation and the two heaviest make up the third generation. Each generation contains of two particles. In addition each particle has its own antiparticle, which shares most its properties but has the opposite charge. Antiparticles were already predicted by Dirac in 1927 [6], as a consequence of the Dirac equation having both positive and negative energy solutions. In 1931 the first antiparticle - the positron - was found in cosmic rays by Anderson. [7]
Quarks interact both strongly, electromagnetic and weakly. Each generation contains one up-type quark and one down-type. Up-type quarks (up, charm and top) have an electric charge of $\frac{2}{3}$, while down-type quarks (down, strange and bottom) have an electric charge of $-\frac{1}{3}$. Each quark exists in three different colours - red, blue and green - and the antiquarks exist in the corresponding anticolours. These colours represent the colour charge, which couples to the strong force just as the electric charge couples to the electromagnetic force. Quarks do not exist as free particles in nature, but always as colourless systems containing either two or three quarks denoted as mesons and baryons, respectively. These systems are commonly known as hadrons. The most "famous" hadrons are protons and neutrons.

Leptons do not interact strongly, but only electromagnetic and weakly. Each generation contains one charged lepton and one neutral. Charged leptons count electrons, muons and tau particles, each with an electric charge of -1. Uncharged leptons count three types of neutrinos - one neutrino per charged lepton. While charged leptons are massive, neutrinos have a mass so small, that it is still unmeasured. However, the mass of the neutrinos should be non-zero in order to explain neutrino oscillations, which describe the transformation between the neutrino types. These oscillations can only occur, if the three types of neutrinos have different masses. Because the neutrino oscillations only depend on the mass difference between the types, the absolute mass of each neutrino is not known. However, the heaviest neutrino is expected to have a mass between 0.04 eV and 0.4 eV [4].

```plaintext
<table>
<thead>
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<th>Quarks</th>
<th>Leptons</th>
</tr>
</thead>
<tbody>
<tr>
<td>up</td>
<td>1 e^-</td>
</tr>
<tr>
<td>charm</td>
<td>1/2</td>
</tr>
<tr>
<td>top</td>
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<tr>
<td>down</td>
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<tr>
<td>strange</td>
<td>1/2</td>
</tr>
<tr>
<td>bottom</td>
<td>1/2</td>
</tr>
</tbody>
</table>
```

Figure 2.1: Particles and interactions of the Standard Model. [5]
All massive particles exist both as right-handed and left-handed particles. Right-handed particles are those with a helicity of +1 meaning, that the spin and velocity of the particles are parallel. Left-handed particles, on the other hand, have a helicity of -1 meaning, that the spin and velocity are antiparallel. Since the velocity and thereby the helicity is not Lorentz invariant, the helicity of massive particles are frame-dependent. For massless particles travelling at the speed of light, it is not the case though. Therefore do neutrinos only exist as left-handed particles and antineutrinos as right-handed particles.

The neutrinos were predicted by Fermi in 1930 [9] in order to explain the continuous energy spectrum of the electron emitted in beta-decay, and in the mid-1950s the neutrinos was measured at the Savannah River nuclear reactor [10]. The first lepton was found already in 1897, when J.J. Thompson saw the electron [3]. Following the discovery of protons in 1911 by Rutherford and neutrons in 1932 by Chadwick, the physicists of the mid-1930s thought they knew all subatomic particles of nature. However, in 1937 the muon was discovered - a new particle having such surprising properties that Rabi, when he heard the news of the discovery, quipped: "who ordered that?" The muons was though first wrongly identified as the Yukawa's particle [11] - later found to be the pion.

2.1.2 The gauge bosons

As mentioned, all three fundamental forces described by the Standard Model are mediated by gauge bosons. They are spin-1 particles and obey Bose-Einstein statistics.

The strong interaction, which is responsible for the nucleon to be bound, is mediated by the massless gluons. Gluons carry colour charge, which implies that they only couple to fermions also carrying colours (the quarks). They were seen experimentally in the 1970's as radiation from electron+proton collisions.

The electromagnetic interaction is mediated by the massless photons, which have been known since Einstein explained the photoelectric effect in 1905. Photons couple to electric charge, though they are electrical neutral themself, and do thus couple to all fermions except from the neutrinos.

The weak interaction is mediated by Goldstone bosons. They count the charged $W^\pm$ bosons and neutral $Z$ boson. $W^\pm$ bosons are the ones responsible for radioactive decay. As the strong interaction couples to colour charge and the electromagnetic interactions to electric charge, the weak interaction do not couple to any specific charge. This non-specific charge is sometimes called the weak charge or the flavour charge, but it is referring to something all particles have. Goldstone bosons are the only massive gauge bosons. They were observed at CERN in 1983, and their masses were found to be 80 and 91 GeV respectively.

2.1.3 Feynman diagrams

All physical processes can be illutrated by drawing Feynman diagrams. Some examples are shown in figure 2.2. Moving from left to right illustrate the timelike evolution going forward in time, while the spatial separation is shown vertical. Lines only connected at one end are known as "external" and represent the observable particles. Lines connected at both ends are "internal" and represent particles, which cannot
be observed directly. The straight, wiggly and curly lines specify fermions, massive bosons/photons and gluons respectively. However, deviations from this convention do occur. Lines represent particles if the arrows are pointing from left to right and represent antiparticles if the arrows are pointing from right to left. This illustrates the fact, that antiparticles correspond to particles going backwards in time. The connection points illustrate the interaction vertices, where the fermions and gauge bosons couple. These interaction points are known as the vertices and are associated with a coupling constant describing the coupling strength.

The loop diagram seen in figure 2.2c has a special role. Since it is internal it can always be added to a process, which make the diagram more complex. The simplest Feynman diagrams are known as first order diagrams. The more complex the diagrams for a given process are drawn, the higher order they achieve.

2.1.4 Transition probabilities

The probability for a transition to occur is proportional to the coupling constant associated with the given process. Observing a process with the initial state $|i\rangle$ and the final state $\langle f|$, there might be several possible interaction paths including different coupling constants. To get the total probability for a transition to occur, all interaction paths need to be considered. This can be done by drawing all possible Feynman diagrams describing the process $|i\rangle \rightarrow \langle f|$ and then pay attention to the vertices that appear.

The transition probability of a given process can be found according to Fermi’s Golden Rule:

$$\lambda_{if} = \frac{2\pi}{\hbar} |M_{fi}|^2 d\tau$$

(2.1)

where $M$ is the transition amplitude or the matrix element and $\tau$ is the phase space available or density of the final states [12].

Since $M$ has to include all possible processes transforming $|i\rangle$ into $\langle f|$, it can be found by making a series expansion in all the coupling constants included in the possible Feynman diagrams. This is the so-called Wick expansion. For each diagram the coupling term associated with $M$ is proportional to the square root of the coupling constant for each vertex.

For the electromagnetic process the coupling constant is $\alpha \approx 1/127$ at low energies and for the weak processes it is $g_w \ll 1$. The largest contribution to $M$ is thus made by the lowest order diagrams.
However, the strong process has a coupling constant in the order of \( \alpha_s \simeq 1 \). Thus the most complex diagrams will give the largest contributions to \( M \). Since it is always possible to draw higher order Feynman diagrams, due to loop diagrams such as the one seen in figure 2.2c, it will lead to infinities. How to come around this problem is described later in this chapter.

### 2.2 Quantum Field Theory

All particles and physical processes of the Standard Model are described mathematically by Quantum Field Theory (QFT), which is a combined theory of Quantum Mechanics and Relativistic Mechanics. QFT is a non-Abelian broken gauge theory based on the symmetry group \( SU(3)_C \times SU(2)_L \times U(1)_Y \), where the suffixes \( C, L \) and \( Y \) represent couplings to colour charge, left-handed particles and weak hypercharge respectively. This will be described later in this chapter. QFT describes the strong interaction by Quantum Chromodynamics, the electromagnetic interactions by Quantum Electrodynamics and combines the electromagnetic and weak interaction in the Electroweak Theory.

In Classical Mechanics all elementary particles are described as point-like particles, while in QFT they only exist as fields. Particles with half integer spin, such as fermions, are described by spinor fields \( \Psi \), spin-0 particles, such as pions or kaons, are described by scalar fields \( \phi \) and spin-1 particles, such as the gauge bosons, are described by vector fields \( \varphi \).

#### 2.2.1 The Lagrange formulation

In the Lagrange formulation of Classical Mechanics a system is described locally by the Lagrangian \( L(\dot{q},q) \), which is a function of the spatial coordinates \( q \) and their time derivatives \( \dot{q} \). The equation of motion is derived from the Euler-Langrange equation:

\[
\frac{d}{dt} \left( \frac{\partial L}{\partial \dot{q}_i} \right) = \frac{\partial L}{\partial q_i} \quad (i = 1, 2, 3)
\]  

(2.2)

The description in QFT is very similar, but because of the relativistic extension the space and time coordinates must be treated on equal footing. In QFT a system is described by the Lagrangian density \( \mathcal{L}(\partial_\mu \phi_i, \phi_i) \), which is a function of the fields \( \phi_i \) and their spatial and time derivatives \( \partial_\mu \phi_i \) with \( \mu = 1, 2, 3, 4 \). The Lagrangian density is often referred to as just the Lagrangian. The relativistic version of the Euler-Lagrange equation is found to be:

\[
\partial_\mu \left( \frac{\partial \mathcal{L}}{\partial (\partial_\mu \phi_i)} \right) = \frac{\partial \mathcal{L}}{\partial \phi_i} \quad (i = 1, 2, 3)
\]  

(2.3)

As for the classical case, applying the Lagrangian for a given system to the Euler-Lagrange equation gives the equation of motion. For scalar, spinor and vector fields, the equation of motion are described by the Klein-Gordon, the Dirac and the Procca equation respectively.
CHAPTER 2. THE STANDARD MODEL

2.2.2 Gauge invariance

One of the good features of the Standard Model is, that all interactions are derived from one single principle - the requirement of local gauge invariance. The Lagrangian must be invariant under the local phase transformation:

\[ \phi_i \rightarrow e^{i\theta} \phi_i \] (2.4)

Applying this transformation has some consequences. For instance, it introduces an interaction term to the Lagrangian describing a free particle, the so-called free Lagrangian. The free Lagrangian for a spinor field \( \Psi \) can be used as an example:

\[ L = i(\hbar c) \bar{\Psi} \gamma^\mu \partial_\mu \Psi - (mc^2) \bar{\Psi} \Psi \] (The Dirac equation) (2.5)

The first term describes the kinematics, while the second term includes the mass. This Lagrangian is not invariant under the local phase transformation, but if every derivative \( \partial_\mu \) is replaced by the covariant-derivative \( D_\mu \equiv \partial + iq/(\hbar c)A_\mu \), the invariance is achieved. By introducing this new field \( A_\mu \) to the kinematic term, the full Lagrangian must also include a free term for the new field. Since \( A_\mu \) is a vector field, the extra term should be the Procca equation. Then a problem arises; the mass term of the Procca equation is not invariant under local gauge transformations. As a consequence, the field \( A_\mu \) has to be massless. \( A_\mu \) can thus be identified as the electromagnetic field mediated by massless photons.

By requiring the Lagrangian to be invariant under local gauge transformation, the coupling between fermions and photons are introduced. Similar calculations would introduce massless gluons, but for the massive gauge bosons the Higgs mechanism must be included.

2.2.3 The Higgs mechanism

As described, the strong and electromagnetic interactions work beautifully with the principle of local gauge invariance. However, to conserve the invariance, the Procca equation requires a massless vector field, which is certainly not fulfilled for Goldstone bosons.

A way to achieve massive Goldstone bosons is by spontaneous symmetry breaking and the Higgs mechanism. The Higgs mechanism introduces a scalar field with \( SU(2) \) symmetry, having a non-zero expectation value in the vacuum. Goldstone bosons acquire mass by interacting with the Higgs field at the cost of local gauge invariance. This is what is meant by the spontaneous symmetry breaking; "spontaneous", because the mass of Goldstone bosons are not assigned initially, but are a consequence of the spontaneous interaction with the Higgs field, and "symmetry breaking", because the Higgs field is not locally symmetric.

The Higgs field has four degrees of freedom, three of them are used by Goldstone bosons, while the one left is associated with the Higgs boson. The Higgs boson couples to all other particles including itself with a coupling constant proportional to their mass. This is how fermions acquire their mass.

Until July 2012, there had been no experimental signs of the Higgs boson. Then July 4th 2012, the ATLAS and CMS experiments announced the observation of a new particle with a mass of almost 126 GeV being consistent with the Higgs
boson within 5\(\sigma\) (with 99.999943\% certainly). Later on, in March 2013, after having analysed more data, the observation was confirmed with a confidence of 10\(\sigma\) (with almost 100\% certainly). Since an observation with a confidence of 5\(\sigma\) is accepted as a discovery, the Higgs boson is now a reality. [13]

2.3 Electroweak Unification

In 1967, Glashow, Weinberg and Salam succeeded in unifying Quantum Electrodynamics (QED) and Fermi Theory, which describes the electromagnetic and weak interactions respectively [14]. In fact, it was Glashow’s original aim to unify these two theories into one single theoretical system. It met some complications, though, such as the attempt to combine two theories with different coupling constants. This was solved by introducing a massless mediator of the electromagnetic interaction and a massive mediator of the weak interaction. This introduced a new problem, why did the mass difference between the two mediators exist? This problem was solved by Weinberg and Salam via the introduction of the Higgs mechanism.

The Electroweak Theory (EW) is based on the symmetry group \(SU(2)_L \times U(1)_Y\), where the suffix \(L\) stands for the coupling to left-handed particles and \(Y\) stands for the weak hypercharge.

Interactions associated with \(U(1)_Y\) include the electromagnetic interaction mediated by photons and the weak neutral current interaction mediated by \(Z\) bosons. Both couple to the weak hypercharge with a coupling constant proportional to the electric charge. Since these bosons are electrical neutral, they cannot perform transitions within a doublet and in addition they preserve flavour. Figure 2.2a gives an example of an allowed interaction.

A symmetry being conserved by the neutral current interactions and the strong interaction as well, is the mirror symmetry also known as parity invariance. It states, that the mirror image of a physical process will be just as possible as the physical process itself. Since left-handed particles are just mirror images of right-handed particles, interactions with left-handed and right-handed particles should occur with equally probabilities.

Interactions associated with \(SU(2)_L\) count the weak charged current mediated by \(W^\pm\) bosons. As \(W^\pm\) bosons only couple to left-handed particles, this part of the EW Theory is maximally parity violating. This was shown by the Wu experiment in the 50s [15]. The coupling between leptons and \(W^\pm\) bosons can only take place within a doublet. However, couplings to quarks are not as simple, since they couple both within the generation and between flavours. The different couplings do not occur with equal probabilities, but are strongest within each generation as described by the CKM-matrix:

\[
\begin{pmatrix}
  d' \\
  s' \\
  b'
\end{pmatrix}
= 
\begin{pmatrix}
  V_{ud} & V_{us} & V_{ub} \\
  V_{cd} & V_{cs} & V_{cb} \\
  V_{td} & V_{ts} & V_{tb}
\end{pmatrix}
\begin{pmatrix}
  d \\
  s \\
  b
\end{pmatrix},
\quad |V_{ij}| =
\begin{pmatrix}
  0.9738 & 0.2272 & 0.0040 \\
  0.2271 & 0.9730 & 0.0422 \\
  0.0081 & 0.0416 & 0.9991
\end{pmatrix}
\]

It means for example, that the coupling constant associated to \(d \rightarrow u + W^-\) will
carry a factor of $V_{ud}$. The matrix elements $V_{ij}$ are theoretically given by the trigonometric functions of the Cabbibo angle [4], but their magnitudes are only known from experiments. As seen from the measured values above, interactions within doubles are indeed favourable. Mathematically this is described by the statement, that the weak eigenstates of the quarks are rotated with respect to the mass eigenstates. Figure 2.2b gives an example of interactions mediated by the $W$ boson.

In the EW Theory, the masses of Goldstone bosons are given by their coupling constants and related by $M_W = M_Z \cos \theta_w$, where $\theta_w$ is the Weinberg angle [14]. However, only relations are predicted - not their actual values. The Weinberg angle was the first to be estimated, but with a very high uncertainty. This was in the late 70s using muon neutrino and antineutrino scattering data. The masses were calculated to be $M_W = 82 \pm 2$ GeV and $M_Z = 92 \pm 2$ GeV. In 1983, Goldstone bosons were seen experimentally at CERN and their masses were found to $M_W = 80 \pm 0.029$ GeV and $M_Z = 91.188 \pm 0.002$ GeV. This measurement was much more precise, so from that time the Weinberg angle was calculated from $M_W$ and $M_Z$ instead [4].

### 2.4 Quantum Chromodynamics

Quantum Chromodynamics (QCD) describes the strong interaction mediated by gluons. It is based on the non-abelian symmetry group $SU(3)$ and contains eight independent colour states. Gluons couple to all particles carrying colour charge - including themselves. This leads to gluon-gluon loops as seen in figure 2.3. As mentioned, the colour charge comes in three different colours - red, blue and green. While gluons cannot change flavour, they do change colour. Since colour charge has to be preserved, gluons are not only carrying one colour but two. This is illustrated in figure 2.4.

#### 2.4.1 Divergencies and Renormalisation

In the case of a tree-level Feynman diagram without any loops as seen in figure 2.2a and 2.2b, the momentum of the internal lines are constrained by the conservation of momenta of the external lines. This is not the case for the diagrams containing loops as seen in figure 2.3. Around the loop, the momentum can take all values equally. Therefore, it is necessary to integrate over all values of the 4-momentum, leading to infinities and divergencies. When adding contributions from all possible diagrams, some of the infinities will cancel out, while others will remain. Since the physical
observables are finite, this problem has to be "fixed". It is done by "absorbing" the divergencies into other quantities. This is know as renormalisation. Often it is done by a redefinition of the coupling constant, to an expression depending on the kinematics and containing one extra degree of freedom such as a momentum cutoff. This expression will have the divergence opposite to the divergence of the loop, which leads to a finite result. As a consequence of the renormalisation the coupling constants $\alpha$ are in fact not constant, but will depend on the transferred momentum $Q^2$. When the coupling constant for a given $Q^2$ is know, it is possible extrapolate to other values of $Q^2$ by the renormalisation group equation.

In QED and EW Theory, the coupling constants are fairly low at low $Q$ and do only increase slowly. For all relevant energy scales the coupling constants are low enough for the Wick expansion to be valid.

In QCD the case is different. For energies around 10 GeV and above, the coupling is measured to be well below 1 and is well described. When extrapolating towards higher $Q^2$ the $\alpha_s$ decrease and the Wick expansion converges quicker. This is shown in figure 2.5. The problem appears when extrapolating towards lower $q^2$. At low $Q^2$ a divergence at an energy scale $\Lambda_{QCD} \approx 200$ MeV appears and $\alpha_s$ becomes larger than 1. At this point the Wick expansion breaks down, as higher orders become more important than lower orders and pertubative calculations cannot be performed anymore. Results from Feynman diagrams will then no longer be valid.

![Figure 2.5: The "running" coupling constant as a function of the energy.[16]](image)

![Figure 2.6: The potential energy between quarks as a function of their separation.[17]](image)

### 2.4.2 Running couplings and asymptotic freedom

One of the great triumphs of QCD was the discovery of the coupling constant being dependent on the interaction distance.

It is well-known from electrodynamics, that the effective coupling depends on the interaction distance. Looking at a point charge embedded in a dielectric medium, the effective charge will depend on the distance from where we are observing due to the screening effect. As a result of vacuum polarisation, a similar effect arise in QED. Vacuum polarisation appears when a photon creates an electron-positron pair.
This pair will only exist for a short while according to the Heisenberg uncertainty principle, but while it exists, the pair will contribute to the polarisation. The same effect appears in QCD. The quark-quark-gluon vertices contribute to an effective screening, increasing the coupling constant at small distances. In addition, also gluon-gluon-vertices appears, but these will have the opposite effect of decreasing the coupling constant at small distances. Since the number of gluons in the Standard Model is larger than the number of quarks, the latter effect dominates and the coupling constant decreases at small distances. This is shown in figure 2.6 (keeping in mind, that a strong coupling corresponds to a large potential).

It leads to two important results. First of all the quarks act as free particles at small distances. This is known as asymptotic freedom, and is the behaviour seen inside hadrons. Secondly, it is not possible to observe two separated quarks, since the potential energy would increase infinitely. This is known as the colour confinement [4].

The relation between figure 2.5 and 2.6 should now be clear. At high transferred energy the wavelengths of the mediators are small, so quarks with a small separation distance are involved. Since the potential energy between these quarks is low at small separation distances, the coupling constant is small as well. So at high transferred energy, the coupling constant is small.
Chapter 3

Probing the hadronic structure and examining the $Z$ boson

After introducing the particles and interactions of the Standard Model, the basis is in order to describe the hadronic structure.

Experimentally, the hadronic structure is explored by making collisions between two hadrons and study the outcomings. If all hadrons were elemental it would be expected, that proton+proton collisions would produce particles in the whole phase space similar to Rutherford scattering. In the 1950s it was observed, in proton-proton collisions at 10 GeV, that particles instead were produced collinear with the beam axis. This would suggest, that protons were not elemental. In addition it was seen in electron-proton collisions, that electrons were deflected, indicating the existence of several hard centres inside protons [18]. This observation increased the desire to understand the inner structure of the proton and was a driving factor in developing QCD. It was suggested that the structure of protons were made up by particles, which exchange large momenta by the electromagnetic interaction, but small momenta by the strong interaction. This model was strengthened when QFT was put on a firm theoretical basis. The conclusion was, that hadrons are made of quarks bound by the strong interaction and therefore theoretically described by QCD.

In how much detail the structure is seen by the experiment is strongly dependent of the collision energy. At high energy collisions, particles are probing with a smaller wavelength and as a result, smaller structures are seen. The collision energy is given by the center-of-mass energy of the colliding system:

$$\sqrt{s} = \sqrt{(E_1 + E_2)^2 - (\vec{p}_1 + \vec{p}_2)^2}$$

where $E_1$ and $E_2$ are the energy of the two colliding particles and $\vec{p}_1$ and $\vec{p}_2$ are their 3-momentum. While $E$ and $\vec{p}$ are frame-dependent, $\sqrt{s}$ is Lorentz invariant, which makes it a suitable quantity for comparing different experiments.

In the attempt to reach as high energies as possible, the development of high energy colliders is of huge importance. Compared to fixed targeted experiments, they can achieve a much larger collision energy, since both of the colliding particles are
highly energetic. As modern particle accelerators reach very high energies - $\sqrt{s} = 7$ TeV at the LHC - they are indeed able to probe the inner structure of the hadrons at the energy scale of $\Lambda_{QCD} \approx 200$ MeV.

3.1 The structure of the proton

The proton consists of two up and one down quark hold together by gluons. These three quarks are called valence quarks and possess most of the energy available inside the proton. In addition, valence quarks describe most of the properties of the proton such as spin, charge etc. This structure is seen at low-energy experiments, while at high energy, the structure is seen in much higher detail. In addition to valence quarks, the proton contains also so-called sea quarks. These quarks appear for a short moment, when gluons create virtual quark-antiquark pairs followed by annihilations. Sea quarks only possess a small fraction of the energy available inside the proton, so most of them are light quarks from the first generation. As illustrated in figure 3.1, this structure provides a whole "sea" of quarks and gluons. Particles of this "sea"; valence quarks, sea quarks and gluons, are commonly known as partons.

As described, the momentum $p_0$ of a proton will be carried by partons. The fraction of the momentum carried by a parton with momentum $p$ is denoted $x = p / p_0$. It means, that $x \approx 1/3$ for valence quarks and $x \ll 1/3$ for sea quarks and gluons. The variable $x$ is the so-called Bjorken-$x$ ($x_B$). It has to be distinguished from the Feynman-$x$, given by $x_F = p_L / p_{L,max}$, where $p_L$ is the longitudinal component of the momentum of a given particle and $p_{L,max}$ is the maximal longitudinal component of the momentum the particle could achieve in the given frame.

The structure of the proton is described by the parton density function (PDF). A PDF tells the probability of finding a given parton with a certain momentum fraction $x$ inside a proton - or in principle inside other hadrons as well. Since this distribution is seen differently for high and low energy experiments, a PDF is always evaluated at a given transferred momentum $Q^2$. 

![Figure 3.1: The inner structure of the proton. [19]](image-url)
Due to the non-pertubative nature of partons, PDFs cannot be computed from first principle. Instead they are determined by comparing theoretical PDF-dependent observables to the values measured by the experiments. Since the first PDF was published in the 1970s, the determination of PDFs has gone through several stages as the theoretical and phenomenological understanding of QCD has been improved. In 1983 the first PDFs based on "global fits" was produced. These PDFs are determined by comparing data and theory for a set of different lepton-hadron and hadron-hadron scattering experiments. By using this wide range of data, the information on the PDFs are achieved. Until this point the analyses were performed at leading order (LO) using the lowest pertubative order in the QCD calculations. As the amount of data from high-precision deep-inelastic scattering increased, next-to-leading order (NLO) analyses became mandatory and today next-to-next-to-leading order (NNLO) analyses are performed. Several groups are currently working on PDF fits, but only a handful have been providing regular updates on NNLO global PDFs. Some of these are the CTEQ/CT, MSTW and NNPDF collaborations. 

In principle there exist thirteen independent PDFs for a given hadron; one for each quark, antiquark and for gluons. In practice, however, the charm and heavier quarks are not measured independently. Such a collection of independent PDFs is called a PDF set. The PDF sets released by the most prominent collaborations contains between five and seven PDFs [20].

Figure 3.2: The MSTW 2008 NNLO PDFs at $Q^2=10$ GeV$^2$ and $Q^2=10^4$ GeV$^2$. [20]

Figure 3.2 shows a typical PDF set at two different energy scales including a 68% confidence level. Is easily seen, that valence quarks dominate at $x \approx 1/3$, while sea quarks and gluons dominate at low $x$ - gluons being the most dominant. In addition it is seen, that observation at higher energy makes the small structures (gluons) more "visible".

p+Pb collisions do not only involve protons but also neutrons. However, the internal structure of neutrons is very similar to that of protons. The valence quarks of the neutron count one up and two down quark (instead of two up and one down).
This modifies the PDF at high $x$, while the dominant sea quarks and gluons are unaffected.
When looking at heavy nuclei, the PDFs get modified. These modifications will be described in the following chapter.

3.2 Drell-Yan production

Having described partons as forming the inner structure of protons and neutrons, the next step is to look at partonic interactions.

The process relevant for the analysis in this thesis is the Drell-Yan process shown in figure 3.3. More general, Drell-Yan processes describe the annihilation of a quark and its antiquark forming a photon or a $Z$ boson, which decay subsequential into an opposite charged lepton pair. The process shown in figure 3.3 is to leading order, while next-to-leading order processes would involve EM and QCD corrections. The EM corrections could include radiative return as seen in figure 3.4. This interaction with photons change the energy of the quark pair, making it possible to form the $Z$ mass. QCD corrections would include gluon radiation as well as gluon loops.

![Figure 3.3: The Drell-Yan process describing the production mechanism.](image)

![Figure 3.4: EM correction - Radiative return.](image)

The Drell-Yan process can be broken up into several steps. At first two suitable quarks have to be found inside the hadrons. For the $p+Pb$ collisions, one quark is found in a proton within the proton beam, while the other quark is found from either a proton or a neutron within the Pb beam. This initial partonic state exist at an energy scale of $\Lambda_{QCD} \approx 200$ MeV and cannot be calculated pertubative as $\alpha_s \sim 1$. The second part describes the annihilation of the two quarks forming a $Z$ boson. This occurs at the energy scale $M_Z \gg \Lambda_{QCD}$ where pertubative calculations become valid. At last the $Z$ boson will decay into two opposite charged muons (or electrons), which could be observed by the experiment.

The differential cross-section for Drell-Yan process $p + n \rightarrow Z \rightarrow \mu^+ + \mu^-$ can
be expressed as:

$$d\sigma = \sum_{ij} \int dx_1 dx_2 f_i^p(x_1)f_j^n(x_2)d\sigma_{ij\to\mu^+\mu^-}(x_1,x_2),$$

(3.2)

where the PDFs $f_1$ and $f_2$ represents the probability of obtaining a quark from each of the colliding hadrons, and $\sigma_{ij\to\mu^+\mu^-}$ represents the partonic cross-section.

This process has to be seen in an environment of multiple interactions. When a p+Pb collision takes place, all quarks and gluons get separated. Due to color confinement, they will immediately combine with quark-antiquark pairs - spontaneously created from the vacuum - to form new hadrons. This is known as hadronisation. As a consequence of the hadronisation, the "memory" of the initial state will get lost for strongly interacting particles. The Drell-Yan process has the advantage, that neither the intermediating $Z$ boson nor the final state muons are strongly interacting. The information carried by the final state muons are therefore neither affected by the nuclear medium nor the following hadronisation.

### 3.3 The Z boson

As described, the mediating Z bosons are not measured by the experiment directly, but are reconstructed from the measured final state particles (muons). If the invariant mass of a measured muon and its antiparticle is equal to the Z mass, the two muons (dimuon) will probably be identical to the final state particles of the Drell-Yan process as shown in figure 3.3. The invariant mass of a dimuon system is found as:

$$m_{\mu\mu} = \sqrt{(E_1 + E_2)^2 - (\vec{p}_1 + \vec{p}_2)^2}$$

(3.3)

where $E_i$ is the energy of the i'th muon and $\vec{p}_i$ is its 3-momentum.

Plotting the number of dimuons as a function of the invariant mass will create a spectrum as seen in figure 3.5. Besides the continuum, this spectrum shows a lot of resonances. Resonances such as the $\phi$, $J/\Psi$ and $\Upsilon$ contain particles identified as bound systems of quarks and their antiparticles - in this case systems of strange, charm and bottom quarks, respectively. These flavorless mesons are commonly known as quarkonia. However, the $Z$-peak differs as it is formed by the $Z$ boson.

The shape of a resonance is given by the Breit-Wigner resonance formula:

$$\sigma(s) = \sigma_{\text{max}} \frac{M^2 \Gamma^2/4}{(s - M^2)^2 + \Gamma^2 M^2}$$

(3.4)

This is the relativistic formula describing the cross-section for measuring two particles with a center-of-mass energy $\sqrt{s}$, forming a resonant state with rest-mass $M$ and a lifetime of $\tau = 1/\Gamma$. As seen from the formula the cross-section for creating a resonant state is enhanced, if the center-of-mass energy of the two particles is equal to the restmass of the resonance state. This is why the resonance has its maximum at $m_{\mu\mu} = M$. When this is the case, the resonance particle is said to be on its mass shell, but it could also be "off-shell". Due to Heisenberg's uncertainty principle, the
Z boson could achieve a mass slightly higher or lower than $M_Z = 91$ GeV. In this case it is said to be "off-shell", but the probability of forming such a state would be reduced according to the Breit-Wigner formula. This is why the resonance is not a delta-function but has a width.

The width of a resonance is given by $\Gamma = 1/\tau$ and can be found by using Fermi’s Golden Rule given in chapter 2 (eq. 2.1). As a result, the decay width of the $Z \to f \bar{f}$ is given by:

$$\Gamma(Z \to f \bar{f}) = N_f \frac{\alpha M_Z}{12 \sin^2 \theta_w \cos^2 \theta_w} \left[ (g_V^f)^2 + (g_A^f)^2 \right]$$

(3.5)

where $N_f = 1$ for leptons, $M_Z$ is the mass of the Z boson, $\theta_w$ is the Weinberg angle and $g_V^f$ and $g_A^f$ are the vectorial and axial weak couplings described by the EW Theory [12].

The width given above is the total fermionic width of the Z-peak given by $\Gamma_Z = \sum_f \Gamma_{f \bar{f}}$. Experimentally it is found to be $\Gamma_Z = 2.4952 \pm 0.0023$ GeV [21]. The Z-peak seen in a spectrum like figure 3.5 does not reflect this full width, since it is only constructed from Z bosons decaying to muons. It is therefore important to distinguish between the different decay modes giving the differential decay widths and the related differential cross-sections. Table 3.1 gives an overview of the different decay modes and their differential decay widths.

### 3.4 Kinematics

To be able to reconstruct Z bosons from the measured muons, it is necessary to define a number of kinematic variables. The variables of relevance are the trans-
CHAPTER 3. PROBING THE HADRONIC STRUCTURE AND EXAMINING THE Z BOSON

<table>
<thead>
<tr>
<th>Z decay modes</th>
<th>$\Gamma_i/\Gamma_Z$ [%]</th>
</tr>
</thead>
<tbody>
<tr>
<td>$e^+e^-$</td>
<td>3.363±0.004</td>
</tr>
<tr>
<td>$\mu^+\mu^-$</td>
<td>3.366±0.007</td>
</tr>
<tr>
<td>$\tau^+\tau^-$</td>
<td>3.367±0.008</td>
</tr>
<tr>
<td>$\nu + \bar{\nu}$</td>
<td>20.00 ±0.06</td>
</tr>
<tr>
<td>hadrons</td>
<td>69.91±0.06</td>
</tr>
</tbody>
</table>

Table 3.1: Decay modes [21].

verse momentum $p_T$, pseudorapidity $\eta$ and rapidity $y$. But before turning to these variables a definition of the ATLAS coordinate system will be in place.

3.4.1 The ATLAS coordinate system

Figure 3.6 shows the coordinate system of the ATLAS detector, in where the beams are moving along the center of the detector - one beam towards the A side and the other beam towards the C side. The $z$-axis is parallel to the beams, the $x$-axis is pointing towards the center of the accelerator and the $y$-axis is pointing upwards. $\phi$ is the azimuthal angle and $\theta$ is the polar angle. The two beams collide at the nominal interaction point (IP) at $(x,y,z)=(0,0,0)$. The real interaction point is often displaced a bit with respect to IP, though.

In practice the angle $\theta$ is rarely used directly. Instead the polar angle is measured in terms of the pseudorapidity $\eta$ defined as $\eta = -\ln\left[\tan\left(\theta/2\right)\right]$. Some corresponding $\eta$ and $\theta$ values are shown in figure 3.7. In terms of momentum, $\eta$ can be written as:

$$\eta = \frac{1}{2} \ln \frac{|p| + p_z}{|p| - p_z}$$

(3.6)

At high energies $|p| \approx E$ and $\eta$ is approximal equal to the rapidity, as will be described later in this section. Since the particle production is constant as a function of rapidity in the absence of overall dynamics, as will be described later in this
section, and since the rapidity is Lorentz invariant, $\eta$ is preferred from $\theta$. In addition the calorimeters in ATLAS are segmented into $\eta$ instead of $\theta$ to achieve a uniform occupation level.

### 3.4.2 Transverse momentum

The transverse momentum $p_T$ is given by:

$$p_T = \sqrt{p_x^2 + p_y^2} \quad (3.7)$$

where $p_x^2$ and $p_y^2$ are the momenta of a given particle along the $x$- and $y$-axis, respectively. The tranverse momentum has a very central role, since the beams only carry momentum along the $z$-axis before the collision, implying $p_T = 0$. Measuring the $p_T$ of the outgoing particles is thus a very important kinematic variable.

### 3.4.3 Rapidity

The rapidity is a kinematic variable defined as [22]:

$$dy = \frac{dp_z}{E} \quad (3.8)$$

where $p_z$ is the momentum of a particle along the beam axis and $E$ is the total energy of the particle. The rapidity is the relativistic analogue of the longitudinal velocity, as is easily seen from eq. 3.8 by letting $E \rightarrow m$.

By integrating the above expression using $E^2 = p_z^2 + p_T^2 + m^2$, the relation between the energy and the rapidity of a particle is obtained:

$$E = m_T \cosh y \quad \text{with} \quad m_T^2 = m + p_T^2 \quad (3.9)$$

For particles with $E \gg m$, it can be shown, by rewriting the above expression, that $e^{-y} = \tan(\theta/2)$. This justifies the equality between rapidity and pseudorapidity. Similar to the pseudorapidity, the rapidity can be rewritten in terms of momentum and energy:

$$y = \frac{1}{2} \ln \frac{E + p_z}{E - p_z} \quad (3.10)$$

The rapidity takes the advantage of being additive, as it is Lorentz invariant. The rapidity of the $Z$ boson is therefore just the sum of the rapidity of the two final state muons or equivalent, the sum of the rapidity of the two quarks from the initial state.

It also means, that if two particles collides with different energies, the combined system will be boosted with respect to rapidity. As an example; if two protons with energies $E_1 = 4.000$ TeV and $E_2 = 1.585$ TeV collide, their rapidities will be $y_1 = 4.53$ and $y_2 = -4.06$, respectively. As the rapidity is additive, the combined system will have a rapidity boost of $y = y_1 + y_2 = 0.47$. This number be relevant later on.
The phase space element $d\tau$, that enters the expression for the transition probabilities given in eq. 2.1 or equivalent and the expression for the cross-section $d\sigma = |M|d\tau$, can be expressed in terms of rapidity using eq. 3.8:

$$d\tau = \frac{dp_x dp_y dp_z}{E} = \frac{1}{2} dp_T^2 d\phi dy$$

(3.11)

If $|M|$ varies slowly with respect to the rapidity, $d\sigma/dy$ is independent of the rapidity. This is the case for at least low values of rapidity, which means that in this range the particle production is uniform with respect to rapidity. As a consequence, a rapidity plateau is seen in the rapidity distribution as shown in figure 3.8 [23].

The width of the plateau depends on the center-of-mass energy of the collision and the mass of the produced particles. It can be found as $\Delta y = 2y_{max}$, which in the case of $Z$-production in p+Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV is given by:

$$y = \frac{1}{2} \ln \left( \frac{E + p_z}{E - p_z} \right) = \frac{1}{2} \ln \left( \frac{E + p_z}{E^2 - p_z^2} \right)$$

$$y_{max} = \frac{1}{2} \ln \left( \frac{(\sqrt{s})^2}{M_Z^2} \right) = \ln \frac{\sqrt{s}}{M_Z}$$

$$\Delta y = 2 \cdot \ln \frac{5.02 \text{ TeV}}{91.2 \text{ GeV}} = 8.02$$

(3.12)

This is a number worth to remember when looking at rapidity distributions later on.
Chapter 4

Heavy Ion Collisions

Heavy ion collisions are strongly affected by the fact, that the colliding particles are not all single nucleons, but composite nuclei. In the following a nucleus with atomic number A is denoted A. Though the complexity of p+A collisions is strongly reduced compared to A+A collisions, several effects have to be taken into account. This includes the geometry of the collision and some nuclear modifications.

4.1 Multiplicity

For each p+A collision, the proton will collide with several nucleons within the nucleus A. Therefore, as the proton is a nucleon itself, several nucleon-nucleon collisions occur. These nucleon-nucleon collisions are referred to as binary collisions. The number of interactions and outgoing particles is therefore larger in A+A collisions compared to p+A collisions. The number of outgoing particles is referred to as the multiplicity. However, the multiplicity is often defined as the number of outgoing charged particles and is measured differentially as a function of pseudorapidity $dN_{ch}/d\eta$. Since the number of nucleons within a nucleus increase with the atomic number, the multiplicity will strongly depend on the mass of the colliding nuclei. This relation is seen in figure 4.1, showing the mean multiplicity of shower particles $\langle n_s \rangle$ as a function of A for at different incident beam energies. These results are obtained from p+A target experiments [24].

![Figure 4.1: Mean multiplicity of shower particles $\langle n_s \rangle$. [24]](image-url)
4.2 Geometry of a heavy ion collision

The multiplicity does not only depend on the number of nucleons within the nucleus, but also on the geometry of the collision. This is illustrated in figure 4.2, which shows the center-of-mass system of a p+Pb collision occurring with impact parameter $\vec{b}$. As illustrated by the gray-scaled tube, binary collisions can only occur in a restricted volume given by the impact parameter and the nucleon-nucleon interaction range. The nucleons within this effective interaction volume are the so-called participants. Collisions with a small impact parameter are referred to as central, while collisions with a large impact parameter are referred to as peripheral. Because the number of participants increase with decreasing impact parameter, the number of binary collisions will be larger for central collisions compared to peripheral, and so will the multiplicity.

![Figure 4.2: Schematical overview of the p+Pb geometry.](image)

Another effect illustrated in figure 4.2 is the Lorentz contraction of the high-relativistic nucleus. This is why the nucleus appears as a flat pancake, when is observed from the center-of-mass system.

4.3 Glauber Model and centrality

In the study of heavy ion collisions it is very important to know if a given collision is central or peripheral. Since the radius of a Pb nucleus is $R \approx r_0 A^{1/3} \approx 7$ fm and the precision of the beam position in ATLAS is in the order of $\sim 0.01$ mm [25], direct measurement of the impact parameter $\vec{b}$ and the corresponding number of participants $N_{\text{part}}$ and binary collisions $N_{\text{coll}}$ are not possible. Instead theoretical techniques have been developed to estimate these quantities from the experimental data using a geometric approach. These techniques are generally referred to as "Glauber Models", but are all variations and extensions of the same original Glauber Model. For this reason $\vec{b}$, $N_{\text{part}}$ and $N_{\text{coll}}$ are commonly known as the Glauber parameters [26].

The history of the Glauber Model began in the 1950s, when Roy J. Glauber addressed the problems regarding the new area of high-energy scattering with composite particles. Glauber's work on putting quantum theory of composite particles on a firm basis succeeded, and his predictions were consistent with the experimental data for protons incident on deuterium and heavier nuclei. Especially when high energy collisions of hadrons and heavier nuclei became possible during the 70s, his work was found to have utility. Following Glauber, several people have made huge progress in
developing the Glauber Model. This work has led to two different approaches - the Optical-limit approximation and the Monte Carlo approach.

Both the Optical-limit approximation and the Monte Carlo approach view the collision of two nuclei as individual interactions between the nucleons. They assume the nucleons to move independently within the nucleus, that the size of the nucleus is large compared to the extent of the nucleon-nucleon forces, and that the nucleons will carry enough energy to be undeflected, when the two nuclei pass through each other. In addition it is assumed, that the nucleon-nucleon cross-section is independent of the number of collisions the given nucleon has already undergone. In addition to these assumptions the calculation of the Glauber parameters require some experimental input parameters.

**Inputs to Glauber Calculations**

The two most important experimental input parameters are the nuclear charge density, measured in low-energy electron scattering experiments, and the energy dependent inelastic nucleon-nucleon cross-section $\sigma_{NN}^{inel}$.

The nuclear charge density is normally described by a Fermi distribution:

$$\rho(r) = \rho_0 \frac{1 + w(r/R)^2}{1 + \exp\left(\frac{r-R}{a}\right)}$$

where $\rho_0$ is the nuclear charge density at the center of the nucleus, $R$ is the nuclear radius, $w$ is the deviation from a spherical shape and $a$ is the "skin depth" describing the surface thickness as illustrated in figure 4.3.

![Figure 4.3: Nuclear charge density.](image)

![Figure 4.4: Inelastic cross-section parameterised by Pythia simulations.](image)

The inelastic nucleon-nucleon cross-section $\sigma_{NN}^{inel}$ is used as an experimental input, because the processes at high energy collisions involve low transferred momentum, making perturbative QCD calculations impossible. The energy dependence of $\sigma_{NN}^{inel}$ is shown in figure 4.4, where the inelastic cross-section is parameterised by Pythia simulations. In addition figure 4.4 shows the total and elastic cross-sections measured.
from data. The points marked by stars indicate the cross-sections used for Glauber Monte Carlo calculations at RHIC.

The Optical-limit approximation

The Optical-limit approximation is the main remaining feature of the original Glauber Model, but was reformulated into a more modern language by Bialas et al. in 1976. This approach assumes a continuous nucleon density and does not locate the nucleons at specific coordinates. The Optical-limit approximation describes the incoming wave as a sum of all possible two-nucleon complex phase shifts with the imaginary part related to the nucleon-nucleon scattering cross-section through the optical theorem [27]. The assumption of independent nucleon trajectories makes it possible to find an analytical expression for the cross-section, the number of binary collisions and the participants.

\[ \hat{T}(b) = \int \hat{T}_A(s) \hat{T}_B(s-b) d^2s, \] (4.1)

where \( \hat{T}_A(s) \hat{T}_B(s-b) \) is the probability of finding a nucleon in the gray-scaled tubes in target A and projectile B. This can be interpreted as the effective overlap area for which a specific nucleon in A can interact with a specific nucleon in B. The probability for an inelastic interaction to occur between the two nucleons is then given by \( \hat{T}(b) \sigma_{NN}^{inel} \). The probability of several nucleon-nucleon interactions to occur are found from a binomial distribution, from where the total inelastic cross-section between nucleus A and B is found. This cross-section and the corresponding number of binary collisions and participants are derived as:

\[ \sigma_{A+B}^{inel} = \int_0^\infty 2\pi b db \left\{ 1 - \left[ 1 - \hat{T}_{AB}(b) \sigma_{AB}^{inel} \right] AB \right\} \]

\[ N_{coll}(b) = AB \hat{T}_{AB}(b) \sigma_{AB}^{inel} \]

\[ N_{part}(b) = A \int \hat{T}_A(s) \left\{ 1 - \left[ 1 - \hat{T}_B(s-b) \sigma_{inel}^{inel} \right] B \right\} d^2s \]

\[ + B \int \hat{T}_B(s-b) \left\{ 1 - \left[ 1 - \hat{T}_A(s) \sigma_{inel}^{inel} \right] A \right\} d^2s \]

Figure 4.5: Schematic representation of the Optical Glauber Model geometry. [26]
where A and B are the numbers of nucleons in nucleus A and B, respectively.

**The Monte Carlo approach**

Due to the increase of computer power the Monte Carlo approach is now the most commonly used approach. It has the advantage of simplicity and ability to simulate the preferred physical observables in contrast to the Optical-limit approximation, which only derives the total inelastic cross-section. In addition it is also possible to emulate the cuts applied to data.

The procedure of the Monte Carlo approach is in principle very simple. Two nuclei are simulated within a three-dimensional coordinate system by distributing A nucleons in nucleus A and B nucleons in nucleus B according to the preferred nuclear distribution \( \rho(r) \). The radius of each nucleon is randomly found from the distribution \( 4\pi r^2 \rho(r) \). Then the collision is simulated for a random impact parameter \( b \), which is drawn from the distribution \( d\sigma/db = 2\pi \). In the very simplest form, a nucleon-nucleon interaction takes place if the distance between two nucleons in the transverse plane is given by \( d \leq \sqrt{\frac{\sigma_{NN}}{\pi}} \). Such a Monte Carlo event is illustrated in figure 4.6 for an Au+Au collisions with \( b = 6 \text{ fm} \), shown both along the x- and z-axis.

![Figure 4.6: Glauber Monte Carlo event. [26]](image)

This description is the basis behind all Glauber Models, but of course many refinements have been done. An example would be the Glauber-Gribov model including nuclear shadowing, fluctuations of the inelastic nucleon-nucleon cross-sections etc.

### 4.3.1 Relating Glauber parameters to data

The Glauber parameters are found by mapping the preferred experimental measured distributions to the corresponding distributions obtained from the Glauber calculations. This is done by using centrality groups. The centrality is defined from 0-100% and is divided into several groups; 0-20%, 20-40% etc. The 20% most central events belong to centrality group 0-20%, while the 20% most peripheral events belong to centrality group 80-100%. With knowledge of the Glauber parameters, the Glauber Model calculates the distributions of the preferred observables within each centrality group. These distributions are then compared to the experimental data. From this
comparison each event is connected to a specific centrality group and thereby to a specific set of Glauber parameters. Since each centrality group covers a range of $b$, the Glauber parameters used are just the average number within each centrality group. Figure 4.7 shows the correlations between the Glauber parameters and the observables.

![Figure 4.7: Correlation between Glauber parameters and final state observables.][26]

The exact mapping procedure depends on the experiment and is a whole study of its own. Therefore, the mapping procedure used in the p+Pb collisions at ATLAS will not be described in detail, but it will be touched briefly later in this thesis.

4.3.2 Relating $N_{coll}$ and $N_{part}$

Even though the Glauber Monte Carlo in principle would give both $N_{coll}$ and $N_{part}$, it is in practice not always the case. A relation between the two parameters is therefore useful.

In A+A collisions the average number of collisions per participant scales as the length $l_z \propto N_{part}^{1/3}$ of the interaction volume along the beam direction. The average number of collisions including all participants, that is the total number of binary collisions, is then given by $N_{coll} \propto N_{part} N_{part}^{1/3} = N_{part}^{4/3}$.

For p+A collisions the relation is more simple. Since the proton beam does only contain one participant, which could interact with the remaining $N_{part} - 1$ participants from the nucleus A, the total number of binary collisions is simplified to $N_{coll} = N_{part} - 1$.

4.4 Nuclear phases

Apart from the geometry, heavy ion collisions also stand out in other areas. This includes the formation of highly compressed partonic phases.
4.4.1 Quark Gluon Plasma

In 1975, it was proposed that superdense matter made of asymptotically free quarks, rather than hadrons, could exist within the framework of experimental high-energy physics. This quark soup with a density of $\rho = 6 \cdot 10^{16}$ g/cm$^3$ was already known theoretically from the Standard Model and was expected to exist in the core of neutron stars, exploding black holes and in the very early universe. However, it was not expected to be an accessible study within the “normal” physics. [28]

This proposal was presented at the same time as the area of high-energy heavy ion collisions took its beginning. At a conference at Bear Mountain in 1974, the opportunities of heavy ion collisions were discussed, and this meeting is now often mentioned as the starting point of heavy ion physics. [29]

This highly compressed partonic phase is now known as the Quark Gluon Plasma (QGP) and is created when hadronic matter is heated and compressed into such a high density, that quarks and gluons experience asymptotic freedom. The formation of the QGP is shown in figure 4.8.

![Figure 4.8: Formation of the QGP. [30]](image)

In A+A collisions the QGP is formed because of the high energy, the Lorentz contraction of the highly relativistic particles and the large compression when the two nuclei collide.

Since the formation of the QGP require the large compression of the two nuclei, it is not expected to be formed in p+A collisions.

4.4.2 Color Glass Condensate

The absence of the QGP in p+A collisions makes it possible to study another phase of matter also formed in heavy ion collisions - the Color Glass Condensate (CGC).

Formation of the CGC is related to the increase of the low-x gluon density at high collision energies as described in section 3.1. The partonic density is described by two evolution equations; DGLAP and BFKL. The DGLAP equation is a function of the transferred energy $\log Q^2$ and the BFKL equation is a function of the momentum fraction $\log(1/x)$ (Bjorken-x). For increasing $Q^2$, the DGLAP equation describes an increase of the number of partons for $x \lesssim 0.1$, but they only occupy a volume of $1/Q^2$, so the overall system becomes more dilute. Complementary the BFKL equation also increases the number of partons as $\log(1/x)$ increases, but the partons will be of the same size. This is shown by the phase diagram in figure 4.9.

In addition, colliding highly relativistic particles get Lorentz contracted, when observed from the center-of-mass system.
As a consequence, the low-x partons - mostly gluons - will become so compressed at high energies, that they will experience asymptotic freedom. The strong interaction between the partons will thus be absent and they will only interact weakly.

This phase of matter is called the Color Glass Condensate as:

- Gluons forming the CGC are colored.
- The evolution of gluons inside the nucleus is Lorentz time dilated and therefore seems slow compared to the "natural" time scale - the same behavior seen in glass.
- The quantum states are highly populated just as a Bose-Einstein condensate for instance.

At low $\log Q^2$ and $\log(1/x)$ an increase in $Q$ and $1/x$ will cause a linear increase of gluons. This increase cannot continue forever, eventually the phase space of the gluons will start to overlap. This phase is know as gluon saturation and is illustrated on the phase diagram in figure 4.9.

\begin{figure}[h]
\centering
\includegraphics[width=0.4\textwidth]{fig4.9}
\caption{The partonic phase diagram. [31]}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=0.4\textwidth]{fig4.10}
\caption{The saturation moment. [32]}
\end{figure}

**Gluon Saturation**

When the phase space of the gluons start to overlap, recombination becomes favorable. After this point, the evolution is non-linear, so the number of created gluons depends non-linear on the number of existing gluons.

The criteria for saturation can be examined by looking at the gluon density and the cross-section for gluon recombination.

The density of the gluons within the nucleus is given by:

$$\rho \sim \frac{x f_A(x, Q^2)}{\pi R_A^2}$$

(4.2)

where $f_A$ is the PDF of nucleus A with radius $R_A$.

Since the cross-section for two gluons to recombine is:

$$\sigma_{gg\rightarrow g} \sim \frac{\alpha_s}{Q^2}$$

(4.3)
the recombination will happen if $\rho \sigma_{gg \rightarrow g} \gtrsim 1$, or equivalent when $Q \lesssim Q_s$, where

$$Q_s^2 \sim \frac{\alpha_s f_A(x, Q^2)}{\pi R_A^2} \sim A^{1/3} \frac{1}{x^{0.3}}$$  \hspace{1cm} (4.4)$$

This is an interesting result, since it states the scale of saturation to be $A$-dependent. This is illustrated in figure 4.10.[33]

### 4.5 Nuclear PDFs

Because of nuclear effects, the PDFs of nucleons within a heavy nucleus is modified compared to the PDFs of free protons. Therefore a distinction is made between proton PDFs describing free protons (pPDF) and nuclear PDFs describing the nucleons within heavy nuclei (nPDF).

Only few collaborations are working with nPDFs and among these, different parameterisations of the nuclear modification are used [34][35][36]. In the following, the EPS09 nPDF set [34] obtained by NLO global fits are presented.

The nPDFs are calculated from the pPDF by introducing a multiplicative nuclear modification factor. For a given nucleus $A$, this is done for each parton flavor $i$:

$$f^A_i(x, Q^2) \equiv R^A_i(x, Q^2) f_i(x, Q^2)$$  \hspace{1cm} (4.5)$$

where $f^A_i(x, Q^2)$ is the nPDF, $f_i(x, Q^2)$ is the pPDF and $R^A_i(x, Q^2)$ is the nuclear modification factor.

The modification factor is parameterised by a piecewise fit function seen in figure 4.11. As seen on the figure, the nuclear modification includes several effects, affecting different regions of the momentum fraction $x$ (Bjorken-$x$). The exact shape of the function depends on the fitting parameters $y_0$, $x_a$, $y_a$, $x_e$, $y_e$ and $\beta$. How these parameters influence the shape of the function is seen in the figure except from $\beta$, which is the slope factor in the Fermi-motion part.

The European Muon Collaboration (EMC) at CERN was the first to observe a difference between the structure of free and bound nucleons using data of deep inelastic scattering of muons on deuteron and on iron [37]. The nuclear modification is therefore often called the EMC-effect, though, this name is often used specific for the region around $x = 0.3$.

There exist several models describing the nuclear effect: models based on single nucleons, pion enhancement, multiquark clusters etc. [38]. Some of these models have also been combined, which has led to a plethora of models sharing the same problem of only providing a good description for certain regions of $x$.

Instead of going into details with the different models, a simplified picture of the nuclear effects are given below.

* As seen in figure 4.11, the modification factor is below one at low $x$. It means that the probability of finding a parton at low $x$ is smaller for heavier nuclei than for free protons. This is caused by the effect of shadowing. In the phase diagram in figure 4.9 is was shown, that the partons expand at low $x$, meaning that the partons begin to shadow each other.
• At higher $x$, there is a region where the probability of finding a parton is larger for heavier nuclei than for free protons. This region is dominated by antishadowing, meaning that the partons do not longer shadow each other, but still the number of partons is larger compared to the proton case.

• The region around $x = 0.3$ is dominated by the EMC-effect. This effect is caused by the fact, that quarks in heavier nuclei move through a larger confinement volume and, as the uncertainty principle implies, carry less momentum than quarks in free nucleons.

• The behavior at $x = 1$ is caused by nuclear Fermi motions. The Fermi motion describes the motion of nucleons within the nuclei, a motion many of the partons within the nucleon will carry as well.

For now the data available for nPDF is very limited. It is therefore not possible to measure the nuclear modification for each parton flavor independently. The nuclear modifications are thus only defined for the valence quarks ($R_A^V$), sea quarks ($R_A^S$) and gluons $R_A^G$. The results obtained by the EPS09 NLO global fits using three experimental inputs - deep inelastic $t$+A scattering, Drell-Yan dilepton production in $p$+A collisions and inclusive pion production in $d$+Au and $p$+$p$ collisions at RHIC - are given in figure 4.12. The thick black line indicate the best-fit result, while the additional lines indicate the combined theoretical and experimental errors.

As seen in the figure, these results describe the shadowing, antishadowing, EMC-effect and nuclear Fermi motions very well.

### 4.6 Binary scaling?

In section 4.3 it is described, that the number of outcoming particles scales with the number of binary collisions. This is known as binary scaling. It is the case for most particles created at high energies. At low energies, on the other hand, the incoming particles do not have energy enough to undergo several binary collisions.
The number of outgoing particles scales therefore with the number of participants divided by two. This is the so-called participant scaling.

However, it is not all outcoming particles from high energy collisions, that scales with the number of binary collisions. The number of particles created from partons with low $x$ is reduced because of the shadowing effect, while the number of particles created from partons with $x \sim 0.1$ is enhanced etc. A typical $x$ value of the partons producing a $Z$ boson can be found from [22]:

$$x_{1,2} = \left[ \frac{M}{\sqrt{s}} \right] e^{\pm y}$$

From this, quarks need an momentum fraction of almost $x=0.02$ to produce a $Z$ boson at mid-rapidity.

As seen from figure 4.12, the area of $x \sim 0.02$ is dominated by antishadowing. However, this effect is not significant at $Q^2 = 100$ GeV and will be even less significant at energies of the LHC because the system gets more dilute as described by the DGLAP evoution equation (see figure 4.9). The production of $Z$ bosons is therefore expected to scale with the number of binary collisions. This is precise what has been observed for the $Z$ production in Pb+Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV with the ATLAS detector [39]. This results is shown in figure 4.13 for both the muons and electron channel and for different $p_T^Z$.

In heavy ion collisions, the nuclear modification of the nucleons varies out though the radius of the colliding nuclei. The nuclear modification of nucleons sitting in the center of the nucleus is expected to be much stronger than of nucleons sitting near the surface of the nuclei. If the nuclear modification factor is assumed not to be equal one, meaning that the particle production do not scale with the number of binary collisions. Then, the nuclear modification factor is expected to be centrality
Figure 4.13: Result of $Z$ production in Pb+Pb collisions showing binary scaling. [39]

dependent. This can be found experimental as seen in figure 4.14. This figure shows the nuclear modification factor as a function of $p_T$ for hadron production at RHIC. It is found as the ratio between the number of produced hadrons scaled by the number of binary collisions in Au+Au collisions and the number of produced hadrons p+p collisions. These hadrons are produced by low $x$ partons and the particle production is therefore reduced at central events because of a stronger shadowing effect. [40]
Figure 4.14: RHIC result showing that nuclear effects are strongest at central collisions. [40]
Chapter 5

The LHC and the ATLAS detector

The Large Hadron Collider (LHC) is the largest and most powerful particle accelerator of the world. It is designed to collide protons at energies up to $\sqrt{s} = 14$ TeV at a luminosity of $10^{34}$ cm$^{-2}$s$^{-1}$ and Pb ions at $\sqrt{s_{\text{NN}}} = 5.5$ TeV at a luminosity of $10^{27}$ cm$^{-2}$s$^{-1}$. The LHC is located at CERN (European Organization for Nuclear Research) near Geneva in a huge underground tunnel with a circumference of 27 km. The machine was completed and began its first run in September 2008. Due to a magnet quench causing a violent helium leakage and large damage on the machine, it was under repair for more than a year before restarting in November 2009. The history of the LHC began already in February 1985, when the machines started excavating the tunnel for its predecessor - the Large Electron-Positron Collider (LEP). When LEP was build the tunnel was made wider than required to make room for the LHC later on.

The LHC consist of two beampipes, kept at ultra-high vacuum, with counter-rotating highly energetic beams of protons or heavier nuclei. The particles come in bunches separated by 25 ns, which are accelerated electric fields. The In addition, 1232 dipole magnets are bending the beam and 392 quadrupole magnets ensure that the beam is in focus. Four places the beams cross and the bunches of particles collide at the interaction point. Each interaction point is surrounded by a detector, making it possible to investigate the outgoing particles. The four detectors includes ATLAS, CMS, ALICE and LHCb as seen in figure 5.1. In addition, this figure illustrates the complexity of the LHC with all its pre-accelerators and experiments.

ATLAS and CMS are the two general purpose detectors. They are designed to discover any new physics that should appear at the high energies. The hunt for the Higgs boson has therefore been a high priority. The search for new physics requires detectors to be able to detect and identify all types of particles emerging from the interaction point. To fulfill this purpose they consist of several sub-detectors, which try to cover all angles around the interaction point. The combination of sub-detectors, make it possible to detect all particles except from the neutrino.
The main purpose of LCBb is to investigating the differences between matter and antimatter. This is done by studying CP symmetry violation in the b-quark sector. ALICE is a heavy-ion detector designed to study the physics of strongly interacting matter at extreme energy densities - the quark gluon plasma.

5.1 The ATLAS detector

The most prominent goal of the ATLAS detector is to discover the origin of the weak symmetry breaking. It should thus be able to measure all decay signatures of the Higgs particle in the mass region predicted by the Standard Model, and it should be sensitive to all new physics. Since many of these processes have a very low cross-section, the LHC must run at very high luminosities requiring the ATLAS detector to have a very efficient background rejection.

As seen on the figure 5.2, the ATLAS detector is build like an “onion” with different layers around the interacting point. These layers consist of sub-detectors detecting different kinds of particles. The sub-detectors can be divided into three parts; the inner detector, the calorimeters and the muon spectrometer. By using all sub-detectors, ATLAS is able to detect almost all kinds of particles and to measure their observables. The inner detector measures the track and momentum of charged particles, the calorimeters determine the energies and the muon spectrometer identifies and measures the momentum of the muons. The only particles not measured in the detector are neutrinos, as they do not interact with any of its components. Instead the neutrinos can be estimated by finding the missing transverse energy \( E_{T}^{\text{missing}} = - \sum E_{T}^{\text{detected}} \). To insure all detectable particles are measured, the detector needs to be hermetic, meaning that all angles should be covered. This goal is almost fulfilled leaving only the region at \( |\eta| > 4.6 \) uncovered. The coverage at small \( |\eta| \) is obtained in the barrel region, in which all detectors are organized into a cylindrical shell parallel to the beam axis. The coverage at large \( |\eta| \) is obtained
in the end-cap region with detectors arranged perpendicular to the beam axis. This structure is easily seen on figure 5.2.

In addition to the various sub-detectors, a complex magnet system is applied to the ATLAS detector. It consist of two parts; a 2 T solenoidal field applied to the inner detector and toroidal field with a field strenght varying between 2 and 8 T applied to the muon spectrometer. As charged particles will bend according to their charge, momentum and the strenght of the field, the magnet system allows the momentum of the particles to be measured.

A last part of the detector is the low level trigger system, which is used for particle identification.

5.1.1 The inner detector

The inner detector is located just a few centimeter from the beam axis as a cylindrical shell with a diameter of 1 meter and a length of 6.2 meters. The inner detector tracks charged particles within $|\eta| < 2.5$ using three different sub-detectors; the pixel detector, the semiconductor tracker (SCT) and the Transistion Radiation Tracker (TRT). These sub-detectors and their structures are seen in figure 5.3. When an event is recorded, it is possible to reconstruct the path of the charged particles from the 3D map of all the hits in the Pixel Detector, SCT and TRT. From these tracks it is possible to reconstrcut the vertex from which the particle was produced.
The pixel detector

The highest granularity is achieved around the interaction point by the pixel detector. Since it is the sub-detector closest to the interaction point, it has the ability to find short lived particles such as B-hadrons. Especially the innermost layer of the pixel detector have this advantage and is thereby known as the B-layer. The pixel detector consist of about 80 million pixel elements divided on 1456 modules in the barrel region and 288 modules in the end-cap region. These modules are covering an area of 1.7 m$^2$. In the barrel region the elements are arranged in three concentric cylindrical layers as seen in figure 5.3 and in the end-cap region, they are located on three disks. All pixels have a size in R-\(\phi\)\(\times\)z of minimum 50\(\times\)400 \(\mu\)m\(^2\), giving an intrinsic accuracy of 10 \(\mu\)m in (R-\(\phi\)) and 115 \(\mu\)m in R. Each track will typical leave three hits in the pixel detector - one hit per layer.

The semiconductor tracker

The semiconductor tracker (SCT) - also named the silicon strip tracker - forms the middle part of the inner detector and is constructed similar to the pixel detector. However, instead of pixel elements, it consist of silicon strips. It contains four layers of silicon detectors measuring 6.36\(\times\)6.40 cm, each containing 780 readout strips of 80 \(\mu\)m pitch, giving around 6.3 million readout channels.

To obtain the space point coordinates each layer contains two silicon detectors - one later with strips parallel to the beam axis and one layer with strips perpendicular to the \(\phi\)-direction. As a consequence the SCT will contribute with four space points per track, with an intrinsic accuracy is of 17 \(\mu\)m in R-\(\phi\) and 580 \(\mu\)m in z.
The Transition Radiation Tracker

The Transition Radiation Tracker (TRT) is the last and outermost part of the inner detector and is designed as a gas straw tracker with the additional ability to detect transition radiation. It consists of 370,000 drift tubes with a diameter of 4 mm, which are located parallel to the beam axis in the barrel and radial in the end-cap. Each drift tube is filled with a mixture of gas - 70% Xenon - and contains a 30 $\mu$m gold plated W-Re wire along the center. This wire is held at an electric potential of about 1.3 kV. When a charged particle passes through a tube, the gas will become ionised and the ionisation electrons will drift towards the wire. It will lead to further ionisation and will cause a detectable voltage spike.

In addition the TRT takes advantage of transition radiation emitted by highly relativistic particles, as they pass through materials with different refractive indexes. For this purpose the area between the drift tubes are filled with a fibrous material. When highly relativistic particles pass through that fibrous material, they will emit radiation in the X-ray region. Since Xenon is very good absorber in the X-ray region, the path of a highly relativistic particle will thus leave a very strong signal in the wires. The only particles at the LHC being relativistic enough to emit transition radiation are electrons and highly energetic muons. Because of the low mass of electrons, they are more likely to detect than muons.

The maximum length of the straws in the barrel is 144 cm, but to reduce the amount of occupancy each straw is divided into two and the signals are readout in both ends. The straws in the end cap are readout at the outer radius. The space points are measured by drift time measurements and provide a resolution of 170 $\mu$m per straw. The TRT contributes typically with 36 tracking points.

5.1.2 The calorimeters

The purpose of the calorimeters is to measure the energy of charged and neutral particles. The calorimeters include the electromagnetic, the hadronic and the forward calorimeters. The electromagnetic and hadronic calorimeters are both located in the barrel and the end-caps, while the forward calorimeter is only located at the end-cap. They are all seen in figure 5.4. The calorimeters are located outside the central solenoid magnet, so they operate in the absence of the electromagnetic field.

The principle behind the energy measurement is to stop the particle in some high density material - the absorbers - by making them produce a shower of secondary particles. These secondary particles will then cause ionisation. By measuring the ionisation level, the energy of the original particle can be estimated. To obtain precise measurements, it is important to measure all showers. The thickness of the calorimeters are thus given in units of radiation lengths and hadronic interaction lengths, both defined as the mean distance over which a particle looses $1/e$ of its energy.
**The electromagnetic calorimeters**

The electromagnetic calorimeter is the innermost of the calorimeters and covers $|\eta| < 1.475$ in the barrel and $1.375 < |\eta| < 3.2$ in the end-cap. In the barrel it has a depth of 24 radiation lengths and 26 in the end-cap. As seen in figure 5.4 the electromagnetic calorimeters in the barrel consist of two identical parts separated by 4 mm leaving the so-called crack region at $|\eta| < 0.1$. In addition they are segmented into three sections in the radial direction. The electromagnetic calorimeters in the end-cap consist of two sections. The electromagnetic calorimeters measure the energy of electrons and photons using lead as the absorber and liquid Argon for ionisation. Since the hadrons do not interact strongly with the high electron density inside the lead, only electrons and photons create showers ionising the liquid argon. The level of ionisation is read out by copper electrodes. The absorber and electrode plates are organized as an accordion-shaped "sandwich" interspersed by the liquid argon. This geometry is shown in figure 5.5. These "sandwiches" are arranged in the radial direction to ensure a fast readout, and therefore the accordion geometry is required to provide a complete coverage without any azimuthal cracks.

![Figure 5.5: Illustration of a shower within the accordion-shaped electromagnetic calorimeters. [43]](image-url)
The hadronic calorimeters

As the hadrons are not stopped by the electromagnetic calorimeters they continue to the hadronic calorimeters. The hadronic calorimeters are located just outside the electromagnetic calorimeters and have a hadronic interaction length of 11. They counts the tile barrel covering $|\eta|<1.0$, the extended barrel covering $0.8 < |\eta| < 1.7$ and the end-cap, which extend out to $|\eta|=3.2$. The tile barrel and extended barrel are divided azimuthally into 64 modules and three layers radial, while the end-cap consist of two disks. In the tile barrel and extended barrel the hadronic calorimeter is using steel as the absorbers and scintillating tiles to measure the showers initialised at the absorbers. Because of the high radiation level the hadronic calorimeters in the end-cap do not use the tile structure, but use liquid Argon like the electromagnetic calorimeters. However instead of using lead as the absorber, the hadronic calorimeters are using copper in the end-cap.

The forward calorimeters

The forward calorimeters are only located in the end-cap, but very close to the beam, making it possible to measure the energy of both electrons, photons and hadrons down to $3.1 < |\eta| < 4.9$. The forward calorimeters consist of three modules in each end-cap. The first has absorbers made of copper to measure the energy of electrons and photons, while the second and third have absorbers made of tungsten to measure the energy of hadronic showers. Each module consists of a metal matrix with longitudinal channels filled with an electrode structure containing concentric rods and tubes parallel to the beam axis. The ionisation material is liquid argon, which is filled in the gaps between the rods and tubes.

5.1.3 The muon spectrometer

The muon spectrometer is the outermost sub-detector of ATLAS and its purpose is to identify muons; to measure their momentum and to perform triggering. As illustrated in figure 5.6, muons are the only detectable particles that can transverse both the inner detector and the calorimeters without being stopped. They will only lose a small part of their energy - around 3 GeV. This means, that muons with an energy below $\sim 2.5$ GeV will never reach the muon spectrometer. Muons are thus identified as the particles reaching the muon spectrometer. However, despite of the material the particles have to pass through, some background might still appear in the muon spectrometer. This could be mesons having pushed their way through or muons from in-flight decay.

A cut-away view of the muon spectrometer is shown in figure 5.7. In this figure the large superconducting air-core toroids magnets are shown as the yellow elements, while the muon chambers are shown in blue. In the barrel the muon chambers are arranged in three cylindrical layers and in the transition and end-cap region they are located on disks - also in three layers.
The high precision tracking system consist of the Monitored Drift Tubes (MDT’s) and Cathode Strip Chambers (CSC’s), which cover a region of $|\eta| < 2.7$. They are constructed similar to the TRT, just with a larger diameter and without the additional feature of transition radiation detection. However, the CSC’s differs a bit since the drifttubes contain multiwires and the cathodes are segmented into strips. For the region $2 < |\eta| < 2.7$ the CSC’s are with higher granularity to withstand the higher rate of particles.
CHAPTER 5. THE LHC AND THE ATLAS DETECTOR

The trigger system consist of the Resitive Plate Chambers (RPC’s) in the barrel and the Thin Gap Chambers (TGC’s) in the end-cap. They are used mainly as trigger detectors because of their fast readout. While the tracking system covers most of the $\eta$-range, the trigger system only covers $|\eta| < 2.4$. Apart from triggering, the RPC’s and TGC’s provides bunch-crossing identification, well-defined $p_T$ determination and measurements of the coordinates in the direction orthogonal to that determined in by the tracking chambers.

5.1.4 The trigger system

The purpose of the trigger system is to identify interesting events, and thereby decide which events should be stored at the ATLAS storage elements and which ones should not. The LHC is designed to run with a collision rate of around 40 MHz for a bunch spacing of 25 ns. It correspond to an interaction rate of approximately 1 GHz. This carres such a huge amount of information, that it is impossible to stored at the ATLAS storage elements. The trigger system should therefore be able to reduce the incoming interaction rate to $\sim 200$ Hz.

The trigger system at ATLAS consists of three levels - level 1, level 2 and the Event Filter. These are denoted as L1, L2 and EF, respectively, and are illustrated in figure 5.8. At each level events passed from the previous level are evaluated, decisions made on the previous level are refined and further cuts are eventually applied. Because of the large interaction rate, some triggers are applied a given prescale factor. It means that only one event per "the given prescale factor" will pass to the next level.

Since the luminosity varies during an entire run, every run is divided into a number of lumi blocks, containing roughly 2 minutes of data taken. Certain decisions made by the triggers will depend on the luminosity and run conditions, and are therefore evaluated within each lumi block. The information about which triggers are used and which prescales are applied, is stored in so-called ”Trigger Menu” - lumi block for lumi block.

L1 is a hardware system sitting close to the readout electronics. At L1 only a limited amount of data are available. All events are evaluated by information from the muon system and reduced-granularity information from all calorimeters. L1 searches for high transverse-momentum muons, electrons, photons, jets, $\tau$-leptons decaying into hadrons, large missing and large total transverse energy. For each event the L1 trigger defines one or more Regions-of-Interest (RoI’s). The information provided by the RoI’s tells which areas in space are of certain interest. L1 makes a decision within 2.5 $\mu$s and reduces the collision rate from $\sim 40$ MHz to $\sim 75$ kHz, which is passed to the L2.

L2 is a software based trigger. L2 evaluates the RoI’s passed from the L1 and refines the decision using the full granularity and precision within the ROIs. It is designed to reduce the rate to $\sim 3.5$ kHz with an event processing time of 40 ms on average.

The final event selection is done by the EF. Its selection is performed using reconstruction procedures and is thus making the most refined decisions. The EF has an average processing time of four seconds and reduces the rate to 200 Hz. This amount of data is small enough to be stored on the ATLAS storage element.
CHAPTER 5. THE LHC AND THE ATLAS DETECTOR

Figure 5.8: A schematic illustration the the trigger system at ATLAS. [44]

5.1.5 Minimum Bias and muon triggers

The triggers relevant for this thesis are the minimum bias trigger, EF_mbMbts_1_1, and the two muon triggers, L1_MU0 and EF_mu8.

The Minimum Bias Trigger Scintillators (MBTS) are hardware based triggers mounted on the end-caps and covers the region of $2.09 < |\eta| < 3.84$. The MBTS L1 fires on all detected signals and passes the events to the L2, firing on signals with $p_T > 100$ MeV. Basically this is the requirement for a minimum bias event, the most "unbiased" event sample accessible. However, this sample also contains events triggered by random triggers. After L2 the events are passed to the Event Filter, but this level inflicts no further requirements. The minimum bias triggers are distinguished depending on the number and location of hits. The EF_mbMbts_1_1 trigger used in this thesis requires at least one hit on each side. Because of the extremely large number of minimum bias events, the MTBS triggers are heavily prescaled. For p+Pb collisions the MTBS triggers have a prescale factor of about 1000.

Another function performed by the MBTS are timing measurements. With a time resolution of $\sim 2$ ns, the time difference between hits in each side of the MBTS provides a good identification of real events [45][46].

L1_MU0 is the online muon trigger described previously and it triggers on all muons. The EF_mu8 is a muon trigger at the Event Filter and triggers on all muons with $p_T > 8$ GeV. While L1_MU0 is prescaled, EF_mu8 is the lowest unprescaled muon trigger used for p+Pb collisions, meaning it is the muon trigger with the lowest
$p_T$ threshold of all unprescaled muon triggers.

5.1.6 Muon reconstructions algorithms

In the ATLAS offline reconstruction software, there exist two main algorithms of muon reconstruction - MUID and STACO [47]. The MUID algorithm fits a global track to hits in the inner detector and the muon spectrometer. STACO uses another procedure, it fits the tracks separately in the inner detector and in the muon spectrometer, and combines the two tracks afterwards. Both of these algorithms correct for the energy loss in the calorimeters and the dead material in the path of the particles.

In most cases the reconstruction algorithms identify muons in both the inner detector and the muon spectrometer and is able to connect the track in the inner detector with the track in the muon spectrometer - these muons are the so-called combined muons. In certain cases, a muons are only reconstructed in the muon spectrometer - these muons are known as stand alone muons. However, muons are never reconstruction alone in the inner detector. Muons reconstructed in the inner detector is known as segment tagged muons, but they still require some hits in the muon spectrometer [\?].

As a result of the two different algorithms, the data sample contains two muon containers - one with STACO muons and one with MUID muons. These are also known as the first and second chain [48]. As the two algorithms have been running on the same data, most of the muons in the two containers are the same physical muons. However, since the two algorithms use different procedures, some muons will only be found with the STACO algorithm, while other will only be found with the MUID algorithm. It is then possible to increase the number of muons by merging the two containers. It has earlier been done by comparing MUID muons with the STACO muons and clarify which muons that recur in both containers and which do not. This is a procedure, which have been performed by each analysis group who wished to increase the number of measured muons. It is also a procedure, which have have been used in earlier stage of this thesis.

In March 2013, the a merged chain were included to the data samples released by ATLAS, and can be used directly. This chain is referred to as the Third Chain and includes muons reconstructed by a combination of the MUID and STACO algorithms [49].
Chapter 6

Measurement of the Z boson

This chapter presents the measurement of Z bosons. First, the Z bosons are measured within different centrality groups and the centrality distribution is evaluated. Secondly, the differential $Z \rightarrow \mu^+\mu^-$ cross-section is evaluated. To measure the differential cross-section, the efficiency of both the trigger system and the reconstruction algorithms, as well as the detector acceptance are evaluated.

The data used in this analysis was taken with the ATLAS detector at the LHC during the 2013 p+Pb run at $\sqrt{s_{NN}} = 5.02$ TeV.

p+Pb collisions are performed with asymmetric beams. The protons in the proton beam is carrying an energy of 4 TeV, while the nucleons in the Pb beam carry an energy of $\sim 1.6$ TeV. As derived in section 3.4.3 a rapidity boost of 0.47 towards the proton going side is therefore expected.

During the p+Pb run the beam direction was reversed. In the first period (period A) the Pb beam was going towards the A side (positive $z$-axis), and in the second period, period B, the Pb beam was going towards the C side (negative $z$-axis). The orientation of the two sides is shown in figure 3.6.

The analysis is performed data driven partly because the production and validation of the Monte Carlo sample has been delayed. It is therefore not possible to compare the data with simulations. However, the data driven method is very ideal, since systematics from the simulation are avoid. However, a simple Monte Carlo sample is used to estimate the detector acceptance.

6.1 Data samples

The data sample used in this analysis is shown in table 6.1 and counts the full statistic available. Run numbers from 217946 to 218589 belong to period A, while run numbers from 218677 to 219114 belong to period B. This analysis uses the Good Run List made by the Heavy Ion Group at ATLAS. The Good Run List specifies, which lumi blocks are “good”. A lumi block is "bad", if the beams were unstable or for other reasons, which could influence the data quality. Among all data stored, only “good” lumi blocks are analysed. For the run numbers marked with (*) in table 6.1, the entire runs are specified as "bad" and is therefore rejected. Some of the rejected lumi blocks are maybe excluded by reasons, which do not affect the analysis of Z bosons. After further examinations in the future, some of the rejected lumi blocks could therefore be included in the analysis. The rejection of the "bad" lumi
blocks reduce the data sample by $\sim 14\%$.

As mentioned in section 5.1.6, the first data sample released did not include the Third Chain. The data used is therefore the reprocessed version, which includes this chain. The reprocessed data also takes the advance of being calibrated and refined. The analysis has been done using all three chains, but since they show a good agreement in most cases only the Third Chain is presented unless anything else is specified.

<table>
<thead>
<tr>
<th>Run</th>
<th>$\int L dt$ [nb$^{-1}$]</th>
<th>Tag</th>
<th>Run</th>
<th>$\int L dt$ [nb$^{-1}$]</th>
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<td>f514_m1312</td>
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<tr>
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<td>f516_m1312</td>
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<tr>
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<tr>
<td>218035*</td>
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<td>f516_m1312</td>
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<td>218048</td>
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<tr>
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<tr>
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<tr>
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<td>f516_m1312</td>
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<td>f514_m1312</td>
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<td>f516_m1312</td>
<td>219114</td>
<td>0.8032</td>
<td>f516_m1312</td>
</tr>
</tbody>
</table>

Table 6.1: The full 2013 p+Pb runlist with the respective integrated luminosities specified. These luminosities are before the Good Run List is applied. All runs marked with (*) are fully excluded by the Good Run List. Run number 218024 was earlier excluded because of missing streams, but has been reprocessed again - it explains the different name.

As a result of the ATLAS trigger system, the data sample exists in several "versions", the so-called streams. Each stream contains data triggered by some specific triggers setup. For this analysis the Hard Probe stream is mainly used, but in order to calculate the trigger efficiencies, the Minimum Bias stream is used as well.
• The Hard Probe Stream:
The Hard Probe stream contains all events triggered by jets, photons, electrons and muons. In this analysis only the muon triggers are used. The Hard Probe stream contains about 80.1 million events before any rejections.

• The Minimum Bias Stream:
The Minimum Bias stream contains all events triggered by the minimum bias triggers (MBTS), which is described in section 5.1.4, insuring an event distribution as unbiased as possible. Before any rejections the number of events in the Minimum Bias stream are about 150 million.

In the following the Hard Probe stream is used unless something else is specified.

6.1.1 Monte Carlo
To estimate the number of $Z$ bosons being unmeasured because of the detector acceptance, a primitive Monte Carlo sample is used [50]. This sample contains $20000 Z \rightarrow \mu^+\mu^-$ events produced with asymmetric proton beams of 4.000 TeV and 1.583 TeV, respectively and a vertex position of $-5.25\pm 56.5$ mm along the z-axis, which corresponds to data. Smearing of momentum has not been applied, since the sample is only used for detector acceptance. In contrast to data, the Monte Carlo sample covers the full $\eta$-range. This Monte Carlo sample does not reflect p+A collisions, it has not been corrected for the energy loss in the detector and is not applied the reconstruction algorithms used in data. However, all this is not necessary for estimation the acceptance.

The Monte Carlo sample is simulated using Pythia 8 [51] as the Monte Carlo generator and recorded in HepMC format [52].

6.2 Event selection
Events containing $Z$ bosons are selected as events containing two opposite signed muons, which pass some specific criterias. The opposite signed muons are referred to as opposite-sign dimuons and the criterias as the offline cuts. Before turning to the selection of $Z$ bosons, all the events are examined as described below.

6.2.1 Minimum bias triggers
All events are required to fire the MBTS in both end-caps with a time difference less than 10ns. By this, events are ensured to be real physical events. In addition, the centrality groups are defined with respect to events triggered by the MBTS-1_1 trigger. The requirement of the MBTS-1_1 to have fired, is therefore important to be consistent with the centrality definitions. It reduces the data sample by approximal 13%.

6.2.2 Pileup
In general, only muons from the primary vertex are of interest. It is therefore a problem if too many particles from secondary vertexes exist in data. These particles
CHAPTER 6. MEASUREMENT OF THE Z BOSON

Figure 6.1: a) Single, b) Double and c) Central diffractive events and their signatures characterised by jets associated with gaps. [54]

are known as pileup. To avoid pileup, events are rejected if a secondary vertex is associated with a $p_T$ of more than 6 GeV. The rejection of pileup could have been done in a more sophisticated way than just by rejecting the whole event. However, since the data do not contain significant pileup, this simple procedure is chosen. From the pileup rejection, the data sample is only reduced by 1.1%.

6.2.3 Diffractive events

Diffractive events are characterised by the exchange of chargeless pomerons and exist either as single, double or central events [53]. These processes are seen in figure 6.1, together with their respective signatures. In principle, it is possible to form $Z$ bosons by the exchange of pomerons in central diffractive events, but because of the high mass of $Z$ bosons, it is very unlikely. Anyhow the diffractive events are rejected, since the centrality groups are defined only for non-diffractive events. As seen in figure 6.1, the diffractive events are characterised by jets associated with a gap. Events with a rapidity gap of 2.1 are therefore rejected. The rapidity gaps are calculated using a tool developed by the Heavy Ion Group at ATLAS - the same tool, which is applied on the data used for defining the centrality groups. It only reduces the data sample by 0.05%, so do not cause any influence of importance.

After rejecting all events as described above, the integrated luminosity is found to be $\int \mathcal{L} dt = 25.5 \text{ nb}^{-1}$ [55].

6.2.4 Criteria for the selection of $Z$ bosons

For events passing the event selection given above, the single muons and opposite-sign dimuons are examined.

For opposite-sign dimuons to be identified as $Z$ bosons, they must have an invariant mass near 91 GeV. Due to the width of the $Z$ peak a mass range of $66 < m_{\mu\mu} < 116$ GeV is chosen. It could be argued to use a smaller range to exclude background, but since the peak has a Breit-Wigner shape with long tails, a smaller mass range would exclude both background and signal. In addition the mass range
of $66 < m_{\mu\mu} < 116$ GeV is widely used in the literature [39][56].

Another requirement is, that at least one of the muons should fire the muon trigger EF\_mu8, which is described in chapter 5. This muon is said to be matched to the trigger. The EF\_mu8 trigger, with a $p_T$ threshold of 8 GeV, is chosen because it is the lowest unprescaled trigger. Since the invariant mass range around the $Z$ boson contains a low level of background, it is enough only to require one muon to be matched. The thought behind is, that if one muon is a real physical muon matched to the trigger, the second muon would probably also be real if the two muons form the invariant mass of the $Z$ boson.

![Figure 6.2: The $p_T$ of each muon forming the invariant mass of the Z boson without any offline cuts applied (except from the mass requirement). It shows the Z peak very clear at $p_T \approx 45$ GeV.](image)

Since the $Z$ bosons have an invariant mass of $\sim 91$ GeV, the two muons should carry $p_T \sim 45$ GeV each. Of course, one muon could carry more energy and the other could carry less, but in most cases the energy is shared equally. It allows for a $p_T$ cut on each muon. Figure 6.2 shows the energy of each muon forming a dimuon within the chosen mass range without any further cuts applied. As seen, the peak at $p_T \sim 45$ GeV is very clear. The $p_T$ threshold is therefore chosen to be $p_T > 10$ GeV. This threshold is examined later on. The $p_T$ cut is very important for the background rejection, since the level of background is high at low $p_T$.

In addition several other offline cuts are applied. Since the muon trigger system does only cover $|\eta| < 2.4$, this is a natural offline cut $\eta$.

To reject muons associated with jets, all muons are required to be isolated. The isolation cut is chosen in such a way, that the $p_T$ of all other particles located in a cone of $\sqrt{d\eta^2 + d\phi^2} = 0.2$ around the muon should not exceed 15% of the $p_T$ of the muon.
In order to avoid cosmic muons and muons from subsequential decay, the reconstructed track of the selected muons should have a closest approach to the primary vertex of less than 10 mm. It means that $d_0 < 10$ mm and $z_0 \sin \theta < 10$ mm, where $d_0$ and $z_0 \sin \theta$ are the impact parameters.

To insure the muon tracking to be good, a cut regarding the fit of the reconstructed track is made. The "goodness-of-fit" is given by the chi-square value and the number of degrees of freedom obtained from the fit. This cut is chosen to be $\chi^2/n_{dof} < 10$. In addition the muon is required to be combined, meaning that the muon is reconstructed in both the inner detector and the muon spectrometer as described in section 5.1.6.

The choice of cuts and thresholds is made partly from the respective distributions as seen in figure 6.3, and partly from the thresholds used by the Heavy Ion Group at ATLAS and in the literature [39][56]. The distributions in figure 6.3 include all muons forming the invariant mass of the $Z$ boson, but without any other cuts. All offline cuts are summarised in table 6.2.

![Figure 6.3: The distributions of variables used for the offline cuts. The distributions include all muons forming the invariant mass of the $Z$ boson, but without any other cuts.](image-url)
CHAPTER 6. MEASUREMENT OF THE Z BOSON

Offline cuts

<table>
<thead>
<tr>
<th>Offline cuts</th>
</tr>
</thead>
<tbody>
<tr>
<td>for events</td>
</tr>
<tr>
<td>exist on the Good Run List</td>
</tr>
<tr>
<td>time difference between MBTS’s in each end-cap &lt; 10 ns</td>
</tr>
<tr>
<td>$\sum p_T^2_{\text{vertex}} &lt; 6 \text{ GeV}$</td>
</tr>
<tr>
<td>rapidity gap &lt; 2.1</td>
</tr>
<tr>
<td>for muons</td>
</tr>
<tr>
<td>$p_T &gt; 10 \text{ GeV}$</td>
</tr>
<tr>
<td>$\eta &lt; 2.4$</td>
</tr>
<tr>
<td>$d_0, z_0 \sin \theta &lt; 6 \text{ mm}$</td>
</tr>
<tr>
<td>$\chi^2/n_{\text{dof}} &lt; 10$</td>
</tr>
<tr>
<td>should be combined</td>
</tr>
<tr>
<td>$p_T^{\text{cone}20}/p_T^\mu &lt; 0.15$</td>
</tr>
<tr>
<td>for dimuons</td>
</tr>
<tr>
<td>$66 &lt; m_{\mu\mu} &lt; 116$</td>
</tr>
<tr>
<td>opposite-sign dimuon</td>
</tr>
<tr>
<td>at least one muon should match the mu8 trigger</td>
</tr>
</tbody>
</table>

Table 6.2: All offline cuts chosen for the selection of Z bosons.

The actual effect of the offline cuts is seen very clear in figure 6.4, which shows the invariant mass peak with and without the cuts applied. As seen, the offline cuts imply a very strong background rejection.

Figure 6.4: The invariant mass peak with and without offline cuts illustrating the effect of background removal.
6.3 Background estimation

To insure all opposite-sign dimuons to be real physical $Z$ bosons from the Drell-Yan process, the background needs to be estimated. Before turning to the estimation it is though necessary to consider the different sources of background.

6.3.1 Background sources

The background is defined as all opposite-sign dimuons passing the offline cuts, but without being real $Z$ bosons from the Drell-Yan process. The area around the mass of the $Z$ bosons is characterised by the absence of significant background. However, some few sources should be mentioned and are given below. Certain background sources are given with an estimate of their significance. These estimates are from an earlier analysis measuring $Z \rightarrow \mu\mu$ in $p+p$ collisions at $\sqrt{s} = 7$ TeV with the ATLAS detector [56]. The contributions from background are also shown in figure 6.5.

- Fake muons from misidentified mesons or muons from light mesons decaying in flight. The muons decaying from light mesons are negligible at high $p_T$. [57]

- The largest and irreducible background to Drell-Yan production, is electrons forming a $Z$ boson by a radiative return: $e^+e^-\gamma \rightarrow Z \rightarrow \mu^+\mu^-$. 

- $t\bar{t}$ decay. The top decays overwhelmingly to $W$ and $b$ quarks (with a branching ratio of $99.99\pm0.09\%$ [21]). $t\bar{t}$ decays with dimuons in the final state can therefore fake the signal. This background is estimated to be 0.1%.

- $Z \rightarrow \tau\tau$, with each $\tau$ particle decaying into a muon (has a branching ratio of $17.36\pm0.05\%$ [21]). This background is estimated to be 0.07%.

- The background from decays of dibosons ($WW$, $ZZ$, $WZ$) is estimated to be 0.2%.

- High $p_T$ muons from $W^\pm \rightarrow \tau\nu$ or $W^\pm \rightarrow \mu\nu$ can pair with a muon from a QCD jet to fake a $Z \rightarrow \mu\mu$ event. This background is found to be negligible.

- The very large cross-section of QCD dijets implies a background due to the high $p_T$ muons in the jets. The QCD background background is estimated to be 0.4%.

- Cosmic muons could also be a background source. However, from the study of non-colliding bunches, this background is found to be negligible.

Background associated to jets is in principle removed by the isolation cut, while cosmic muons are rejected by the requirement on the closest approach. However, it cannot be expected that the offline cuts exclude all this background and do certainly not exclude background from radiative return, $t\bar{t}$ decay, $Z \rightarrow \tau\tau$ and diboson decay.

It is therefore important to estimate the level of background. It could be done by comparing the data with Monte Carlo simulations, but also datadriven methods of background estimation exist.
The background could be estimated by fitting the invariant mass spectrum in the range of \( m_{\mu\mu} < 66 \text{ GeV} \) and \( m_{\mu\mu} > 116 \text{ GeV} \). These two regions are expected only to contain background. The obtained fit is then extrapolated to the region of \( 66 \leq m_{\mu\mu} \leq 116 \text{ GeV} \). The area below the fitting function in the range of \( 66 \leq m_{\mu\mu} \leq 116 \text{ GeV} \) is then counted as the background. This is known as the side-band method. However, this procedure has a tendency to estimate a too large background because of the high background level at low \( p_T \).

Another choice could be the so-called ABCD-method. The basic idea behind this method is; first to select a number of known background signals, which could be muons from jets. A number of uncorrelated cuts removing this background are then found. By applying the same cuts on the signal, the reduction of signal is an estimate of the background.

In this analysis the background is estimated from the number of same-sign dimuons and is described below. It is the method used by the Heavy Ion Group at ATLAS. \[39\]

### 6.3.2 Background estimation from same-sign dimuons

Dimuon production from jets do not prefer opposite-sign dimuons from same-sign, so the number of opposite-sign dimuons is expected to be equal the number of same-sign dimuons. The QCD background can therefore be estimated as the number of same-sign dimuons. On the other hand, this estimation will not include background from processes forming opposite-sign dimuons such as radiative return (irreducible), \( t\bar{t} \)-decay, \( Z \rightarrow \tau\tau \) and diboson-decay. Since the background level is very low and the QCD background is the dominant source of the reducible background, as described in the previous and as seen in figure 6.5, other background sources than the QCD are
expected to be almost negligible. Further study of these backgrounds are therefore not performed in this analysis.

Table 6.3 shows the number of opposite-sign dimuons, same-sign dimuons and the background to signal ratio for the three chains. The "Background" is the number of same-sign dimuons, while the "Signal" is the number of opposite-sign dimuons subtracted the background. As seen in the table, the estimated background of the Third Chain and of muons from the STACO algorithm are consistent, within the statistical errors, with the result of 0.4% obtained by [56]. The muons from the MUID algorithms is shown to have a smaller background, but it is not significant. It is also seen, that the goal of increasing the number of muons with the Third Chain is achieved.

<table>
<thead>
<tr>
<th>Chain</th>
<th>N_{opposite-sign}</th>
<th>N_{same-sign}</th>
<th>bkg/sig [%]</th>
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</thead>
<tbody>
<tr>
<td>Third Chain</td>
<td>1941</td>
<td>6</td>
<td>0.3101 ± 0.1336</td>
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<tr>
<td>MUID</td>
<td>1918</td>
<td>5</td>
<td>0.2624 ± 0.1229</td>
</tr>
<tr>
<td>STACO</td>
<td>1828</td>
<td>6</td>
<td>0.3293 ± 0.1430</td>
</tr>
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</table>

Table 6.3: Background estimation for the three chains using the number of same-sign dimuons.

Figure 6.6 shows the invariant mass peak of the $Z$ boson for both opposite-sign and same-sign dimuons. Despite of the very low background, it is seen that the QCD background primarily exist at low invariant mass as expected (see figure 6.5).

Figure 6.6: Invariant mass distribution for opposite-sign and same-sign dimuons.
6.4 Optimisation of the offline cuts

Having introduced the method of background estimation, it is now possible to evaluate the offline cuts given in table 6.2. Evaluating a given cut is done by calculating the background to signal ratio for different thresholds. The threshold given the best background to signal ratio is then chosen. The threshold could have been evaluated by other parameters, but the background to signal ratio is chosen inspired from an earlier measurement of Z bosons in Pb+Pb collisions [57]. Among all cuts applied, especially the $p_T$ and isolation cut are important regarding the background rejection. These two cuts are therefore evaluated keeping other cuts constant as given in table 6.2. The isolation threshold is evaluated keeping the $p_T$ threshold constant at $p_T > 10$ GeV and the $p_T$ threshold is evaluated using keeping the isolation threshold constant at $p_T^{cone20}/p_T^{\mu} < 0.15$. The background to signal ratios for different thresholds are seen in table 6.4 and 6.5.

<table>
<thead>
<tr>
<th>Threshold</th>
<th>$N_{\text{opposite-sign}}$</th>
<th>$N_{\text{same-sign}}$</th>
<th>bkg/sig [%]</th>
</tr>
</thead>
<tbody>
<tr>
<td>no threshold</td>
<td>2029</td>
<td>51</td>
<td>2.6</td>
</tr>
<tr>
<td>0.30</td>
<td>1957</td>
<td>9</td>
<td>0.46</td>
</tr>
<tr>
<td>0.20</td>
<td>1948</td>
<td>7</td>
<td>0.36</td>
</tr>
<tr>
<td>0.15</td>
<td>1941</td>
<td>6</td>
<td>0.31</td>
</tr>
<tr>
<td>0.10</td>
<td>1907</td>
<td>6</td>
<td>0.32</td>
</tr>
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</table>

Table 6.4: The background to signal ratio for different isolation thresholds.

<table>
<thead>
<tr>
<th>Threshold for $\mu_1$</th>
<th>Threshold for $\mu_2$</th>
<th>$N_{\text{opposite-sign}}$</th>
<th>$N_{\text{same-sign}}$</th>
<th>bkg/sig [%]</th>
</tr>
</thead>
<tbody>
<tr>
<td>10</td>
<td>10</td>
<td>1941</td>
<td>6</td>
<td>0.31</td>
</tr>
<tr>
<td>10</td>
<td>8</td>
<td>1960</td>
<td>9</td>
<td>0.46</td>
</tr>
<tr>
<td>8</td>
<td>8</td>
<td>1966</td>
<td>9</td>
<td>0.46</td>
</tr>
</tbody>
</table>

Table 6.5: The background to signal ratio for different $p_T$ thresholds.

As seen in the tables, the background is in general low for all thresholds. However, by removing the isolation requirement, the QCD background is increased significantly, since it allows muons from jets to form dimuons. Apart from that, the effect of varying the cuts is small. From this variation, the thresholds are chosen to be $p_T > 10$ GeV and $p_T^{cone20}/p_T^{\mu} < 0.15$ as previously stated in table 6.2.

6.5 Background corrected plots

Since the collisions are expected to have a rapidity boost of 0.47 in the proton direction, the two periods need to be corrected with respect to the geometric change caused by the reversed beam directions. This is done by switching the sign of $\eta$ and rapidity in period B, but it will need some validation. Figure 6.7 shows the rapidity distribution of the Z bosons with reversed rapidity for period B.

It shows an overall good agreement. Running a Kolmogorov Test [58] gives $D = 0.989$, meaning that the two distributions arise from the same distribution.
Figure 6.7: Rapidity distribution of the Z bosons shown separately for period A and for period B. As seen they show a good agreement.

with a probability of 98.9%. This small comparison acts as a simple validation of combining the periods, as will be done in the following.

Because of the asymmetric beams, a rapidity boost of 0.47 is expected. As seen in the rapidity distribution, this boost is not present, the mean rapidity of both period A and B are compatible with zero. A reason for this could be the large rapidity plateau with a width of 8.02. Since the acceptance of the detector only covers a rapidity range of 5.4, we might only see the central part of the plateau. As a consequence, a displacement 0.47 will not be observed. To make a cross-check, corrections will need to be applied and comparison with Monte Carlo should be done. Since the Monte Carlo is not yet released, it is not done in this analysis.

6.6 Centrality dependence

6.6.1 Centrality in the data

As described in chapter 4, the centrality for a given event is related to the multiplicity, but the actual physical quantities defining the centrality groups differs between experiments. For p+Pb collisions the centrality groups are defined through the transverse energy deposit in one of the forward calorimeters. The reason for using the transverse energy deposit in the forward calorimeters instead of the electromagnetic and hadronic calorimeters is to separate quantities defining the centrality from the quantities used in the centrality analysis.

Because of the asymmetric collisions, only the calorimeter in the Pb-going side is used. It has been shown in the p+Pb pilot run that the transverse energy in that direction appears to be mostly directly correlated with the p+Pb collision geometry
The centrality groups are defined in the minimum bias data for all events triggered by the MBTS L1 trigger, which has a reconstructed vertex, good MBTS timing, at least two tracks with $p_T > 100$ MeV and should satisfy the offline tracking conditions [59]. The definitions of the centrality groups are shown in table 6.6 and the $\sum E_T$ distribution of the total number of minimum bias events is shown in figure 6.8 for six different centrality groups.

<table>
<thead>
<tr>
<th>Centrality</th>
<th>FCal $E_T$ range [GeV]</th>
<th>$N_{\text{part}}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>0-1%</td>
<td>$E_T &gt; 90.8727$</td>
<td>$18.26^{+3.24}_{-1.91}$</td>
</tr>
<tr>
<td>1-5%</td>
<td>$90.8727 &gt; E_T &gt; 65.839$</td>
<td>$16.59^{+2.04}_{-0.92}$</td>
</tr>
<tr>
<td>5-10%</td>
<td>$65.839 &gt; E_T &gt; 53.6545$</td>
<td>$15.59^{+1.62}_{-0.87}$</td>
</tr>
<tr>
<td>10-20%</td>
<td>$53.6545 &gt; E_T &gt; 39.9624$</td>
<td>$13.05^{+1.04}_{-0.73}$</td>
</tr>
<tr>
<td>20-30%</td>
<td>$39.9624 &gt; E_T &gt; 31.0642$</td>
<td>$11.36^{+1.64}_{-0.64}$</td>
</tr>
<tr>
<td>30-40%</td>
<td>$31.0642 &gt; E_T &gt; 24.1534$</td>
<td>$9.79^{+0.59}_{-0.61}$</td>
</tr>
<tr>
<td>40-60%</td>
<td>$24.1534 &gt; E_T &gt; 13.5539$</td>
<td>$7.42^{+0.49}_{-0.59}$</td>
</tr>
<tr>
<td>60-90%</td>
<td>$13.5539 &gt; E_T &gt; 2.81714$</td>
<td>$4.06^{+0.21}_{-0.37}$</td>
</tr>
</tbody>
</table>

Table 6.6: Centrality definitions in data. [55]

Figure 6.8: $\sum E_T$ distribution with the six centrality groups indicated. [60]

### 6.6.2 Correction of FCal $E_T$

The energy deposit in the forward calorimeters and the vertex position are found to differ between period A and B. The reason for this is not well understood, but for now it is bypassed by correcting the FCal $E_T$ for period B. The correction is done by implementing a common tool made by members of the Heavy Ion group at ATLAS. This tool adds a correction to the FCal $E_T$ for period B depending on the energy deposit and the vertex position. Figure 6.9 shows the FCal $E_T$ distribution of period A with respect to period B both with and without the correction. As seen in the figure, the uncorrected FCal $E_T$ deviates between the two periods, while the corrected FCal $E_T$ seems to be consistent. Figure 6.10 shows the distribution of the
vertex positions in the two periods, which again shows a difference between the two periods. The mean vertex position for period A is $-3.05 \pm 0.03$ and is $-3.94 \pm 0.03$ for period B. The two distributions are therefore not even consistent within the statistical errors.

Figure 6.9: Ratio between FCal $E_T$ in period A and B with an without the correction applied.

Figure 6.10: The vertex position in period A and B showing their means not to be consistent.

### 6.6.3 Centrality plots

Table 6.7 gives the number of $Z$ bosons (background subtracted), the number of same-sign dimuons, the total number of events triggered by the MBTS [55] and the number of binary collisions measured as $N_{\text{part}} - 1$ - all within each centrality group.

<table>
<thead>
<tr>
<th>Centrality group</th>
<th>$N_Z$</th>
<th>$N_{\text{same-sign}}$</th>
<th>$N_{\text{events}}$</th>
<th>$N_{\text{coll}}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>0-1%</td>
<td>75</td>
<td>0</td>
<td>5.44347e+08</td>
<td>17.26$^{+3.24}_{-1.01}$</td>
</tr>
<tr>
<td>1-5%</td>
<td>203</td>
<td>0</td>
<td>2.17341e+09</td>
<td>15.55$^{+2.04}_{-0.92}$</td>
</tr>
<tr>
<td>5-10%</td>
<td>224</td>
<td>0</td>
<td>2.73261e+09</td>
<td>14.59$^{+1.62}_{-0.87}$</td>
</tr>
<tr>
<td>10-20%</td>
<td>364</td>
<td>1</td>
<td>5.46006e+09</td>
<td>12.05$^{+1.84}_{-0.73}$</td>
</tr>
<tr>
<td>20-30%</td>
<td>284</td>
<td>3</td>
<td>5.44568e+09</td>
<td>10.36$^{+1.67}_{-0.64}$</td>
</tr>
<tr>
<td>30-40%</td>
<td>255</td>
<td>1</td>
<td>5.46372e+09</td>
<td>8.79$^{+0.61}_{-0.59}$</td>
</tr>
<tr>
<td>40-60%</td>
<td>322</td>
<td>1</td>
<td>1.09220e+10</td>
<td>6.42$^{+0.49}_{-0.59}$</td>
</tr>
<tr>
<td>60-90%</td>
<td>189</td>
<td>0</td>
<td>1.63799e+10</td>
<td>3.06$^{+0.21}_{-0.37}$</td>
</tr>
</tbody>
</table>

Table 6.7: The number of measured $Z$ bosons (background subtracted), background, the total number of events triggered by the MBTS and the number of binary collisions within each centrality group.

As described in section 4.6 it is interesting whether binary scaling is observed or not. It gives an indication of, in which degree the initial state quarks are affected.
by shadowing or antishadowing. A useful quantity to illustrate is the number of $Z$ bosons within each centrality group scaled by the number of binary collisions. As seen from table 6.7 the centrality groups do not have the same width with respect to centrality. The number of $Z$ bosons are therefore divided by the total number of minimum bias events within the same centrality group. Since the centrality is defined linear with respect to the total number of minimum bias events, the $N^Z / N_{\text{events}}$-ratio is unaffected by the width of the centrality groups. The centrality distribution is seen in figure 6.11.

![Figure 6.11](image-url)

Figure 6.11: The centrality distribution of the $Z$ bosons scaled by the number of binary collisions. The gray-scaled squares gives the systematic uncertainties on the number of binary collisions.

Figure 6.11 shows, that the number of measured $Z$ bosons scaled by the number of binary collisions increase strongly as the collisions get more central. The binary scaling is therefore not seen, since this would be described by a line with a slope of zero as seen in earlier Pb+Pb measurements [39] presented in section 4.6. As described in section 4.6, deviation from binary scaling could occur if the $Z$ bosons are produced by partons affected by shadowing or antishadowing. The result shown in figure 6.11 could indicate an effect of antishadowing, which would lead to a larger production of $Z$ bosons in central events. However, this effect is not expected and not earlier observed as described in section 4.6.

Another possibility, causing the deviation from binary scaling, could be, that the Glauber Model is not sufficient. The default Glauber Model used rely only on geometric considerations without any corrections for nuclear modifications or fluctuations of the nucleon-nucleon cross-sections. Figure 6.12 and 6.13 show the same centrality distribution, but instead of using the default Glauber Model, the Glauber-Gribov Model is used to measure the Glauber parameters. It is beyond the scope of this thesis to go into details with the Glauber-Gribov model [61], but the
effect on the two plots is caused by the fluctuations in the nucleon-nucleon cross-section [62]. The Glauber-Gribov Omegas of 0.5 and 1.0 used in the two plots describe the degree of fluctuations - \( \Omega = 1.0 \) describes the strongest fluctuation of the two. By comparing the slope of all three plots, the fluctuations of the nucleon-nucleon cross-section seems to be necessary if the \( Z \) production should scale with the number of binary collisions. The Glauber parameters from the Glauber-Gribov model are measured as test trial by the Heavy Ion Group at ATLAS and has been given without systematic uncertainties. This is why figure 6.12 and 6.13 are shown without.

![Figure 6.12: Centrality distribution of \( Z \) production using the Glauber-Gribov model with \( \Omega=0.5 \).](image1)

![Figure 6.13: Centrality distribution of \( Z \) production using the Glauber-Gribov model with \( \Omega=1 \).](image2)

6.7 Efficiencies and acceptance

Despite of the intension to measure all outcoming particles, some \( Z \) bosons might remain unmeasured. This is caused by inefficiencies in the trigger system and the reconstruction algorithms and is caused by the acceptance of the ATLAS detector. In the following, an estimation of the efficiencies is performed.

6.7.1 Trigger efficiencies

The trigger efficiency of the \text{mu8} trigger - the one used in this analysis - is first measured for single muons and later transformed into the efficiency of measuring \( Z \) bosons. It is the procedure used in [63].

As described in chapter 5, the L1 triggers are hardware based triggers and a part of the detector. It means, that the trigger efficiency is \( \eta \) and \( p_T \) dependent, and it is therefore measured as a function of these.
CHAPTER 6. MEASUREMENT OF THE Z BOSON

Single muon trigger efficiencies

As a first guess, the efficiency of the mu8 trigger could be measured by dividing the number of muons firing the trigger by the total number of muons. Unfortunately it is not that simple. Since the Hard Probe stream only contains events triggered by very specific triggers, that "total number of muons" would be strongly biased. Instead of using only the Hard Probe stream, it is necessary to look at the Minimum Bias stream, since it contains the most unbiased data available. Now the problem is, that the Minimum Bias stream is so strongly prescaled, that it does only contain very few high $p_T$ muons. This problem is solved by including a supporting trigger, the MU0 trigger at L1. The L1 MU0 trigger fires in principle on all muons as described in chapter 5. It is then possible to calculate the efficiency of the MU0 trigger in the Minimum Bias stream and to calculate the efficiency of the mu8 trigger in the Hard Probe stream with respect to the MU0 trigger. The real efficiency of mu8 is then found as the product of the two aforementioned efficiencies:

$$
\epsilon_{\mu\text{trig - mu8}} = \epsilon_{\mu\text{trig - mu8 wrt MU0}} \cdot \epsilon_{\mu\text{trig - MU0}}
$$

To be able to measure the efficiency of the MU0 trigger, no $p_T$ threshold is applied on the muons. The trigger efficiency is therefore not done using muons forming the invariant mass of the Z boson, but counts all muons. To ensure the muons to be "good" physical muons, all other offline cuts listed in table 6.2 are applied.

Figure 6.14a and 6.14b show the trigger efficiency of the mu8 trigger with respect to MU0. Figure 6.14a clearly shows a large drop at 8 GeV, which is expected because of the trigger threshold. It also shows that the efficiency is a bit lower at 8 GeV compared to at higher $p_T$. This is a tendency all triggers show just above their thresholds and it is the reason why the mu8 trigger is used instead of the MU10 - simply to ensure the efficiency to be as high as possible at $p_T = 10$ GeV. At 10 GeV and above the efficiency is taken to be constant as a function of $p_T$.

Figure 6.14b shows the trigger efficiency of the mu8 trigger with respect to the MU0 as a function of $\eta$ for all muons with $p_T > 10$ GeV. As seen in this plot, the efficiency is clearly $\eta$ dependent and reflects the different regions in the detector as illustrated in the figure.

The result obtained by calculating the trigger efficiencies for all three chains is quite striking. As seen in figure 6.14c the efficiency obtained using STACO muons is very good, while both MUID and Third Chain muons give a lower efficiency at mid-rapidity. Why this is so, is not well understood, but will need further investigations in the future.

Since the beams are asymmetric, it seems reasonable if the trigger efficiency is period dependent. It would be caused by higher occupancy in the trigger in the
proton going direction, but as seen in figure 6.14d the trigger efficiency is not period dependent.

The average efficiency of the mu8 trigger with respect to the MU0 is found to be:

\[ \epsilon_{\text{trig - mu8 wrt MU0}} = 0.975^{+0.000188}_{-0.000189} \text{ for muons with } p_T > 10 \text{ GeV} \]

This high efficiency is very reasonable since the mu8 trigger is expected to fire on all muons with \( p_T > 8 \) GeV having already fired the MU0 trigger.

Figure 6.14: Trigger efficiencies of the mu8 trigger with respect to the MU0 trigger.

Figure 6.15a and 6.15b show the trigger efficiency of the MU0 trigger measured in the minimum bias. The heavy prescale applied to the Minimum Bias stream is clearly shown by the low statistic at high \( p_T \). Also the efficiency of the MU0 trigger shows a large drop at low \( p_T \), but above 10 GeV, the efficiency has reached its maximum. As seen in the figure, the efficiency of MU0 is quite low. It has an average efficiency of:

\[ \epsilon_{\text{trig - mu8}} = 0.842^{+0.0134}_{-0.0140} \text{ for muons with } p_T > 10 \text{ GeV} \]

However, this is expected from earlier trigger performance analysis [64]. The \( \eta \)-dependence in figure 6.15b also shows a plateau at mid-rapidity with low efficiency.
This is consistent with earlier analysis of L1 triggers. Measurements of the L1 MU11 has shown an efficiency of 0.725 in the barrel region and 0.935 in the end-cap region [64]. Apart from the plateau of relative low efficiency in the barrel region a dip is seen at $\eta \approx 0.8$. When seperating the two periods, the dip at $\eta \approx 0.8$ is shown to be from period A as seen in figure 6.15c. A physical reason for this, such as detector segments out of function, do not seem reasonable since all events exist on the Good Run List. In addition the efficiency has been investigated run by run and it has been found, that the dip at $\eta \approx 0.8$ is not limited to few successive runs. The dip must therefore just be a result of statistical fluctuations. As seen in figure 6.15d, the efficiency of the MU0 trigger using the three different chain show consistency.

As mentioned, the total single muon efficiency of the mu8 trigger is given as the product of the two efficiencies measured above. The combined result is shown in figure 6.17.

The total single muon trigger efficiency for muons with $p_T > 10$ GeV is measured to be:

$$\epsilon_{\text{trig - mu8}} = 0.820^{+0.0132}_{-0.0138} \text{ for muons with } p_T > 10 \text{ GeV}$$
CHAPTER 6. MEASUREMENT OF THE Z BOSON

Efficiency of the Z bosons

Since the trigger efficiencies is shown to be independent of \( p_T \) for all muons above the \( p_T \) threshold, the final trigger efficiency of the Z bosons takes only the \( \eta \) dependence into account.

The trigger efficiency of the Z bosons depends on whether both muons have fired the trigger or just one of them. For all Z bosons with only one muon matched to the trigger, the trigger efficiency of the Z bosons is equal to the single muon efficiency. The efficiency of Z bosons with two matched muons is given by:

\[
\epsilon_{\text{trig}}^Z = 1 - (1 - \epsilon_{\text{trig}}^\mu(\eta_1))(1 - \epsilon_{\text{trig}}^\mu(\eta_2))
\]

Among all measured Z bosons, 506 includes only one matched muon and 1435 includes two matched muon. It leads to a total trigger efficiency of the Z bosons of:

\[
\epsilon_{\text{trig}}^Z = 0.926^{+0.0291}_{-0.0304}.
\]

The trigger efficiency is later applied to the rapidity distribution. The final trigger efficiency is therefore found as a function of rapidity. This is shown in figure 6.17.

6.7.2 Reconstruction efficiencies

Another reason for unmeasured Z bosons is the ineffectivity of the reconstruction algorithms. Since all muons forming the Z bosons are required to be combined, the total reconstruction efficiency is the product of three efficiencies: 1) the reconstruction in the inner detector, 2) the reconstruction in the muon spectrometer and 3) the efficiency of combining the two. All these efficiencies have to be taken into account.

The reconstruction efficiencies is found by the tag and probe method using \( Z \to \mu\mu \) events. The idea behind this method is first to find a muon, which pass some strong criterias - this is the tag. Then another muon, which forms the invariant mass of the Z boson and has the opposite sign of the tag, is found - this is the probe. From examination of the probe, the reconstruction efficiency is found.
Efficiency for the inner detector

The reconstruction efficiency of the inner detector is found as the fraction of probes being reconstructed in the inner detector with respect to all reconstructed muons. As described in section 5.1.6, all combined muons are reconstructed in both the inner detector and muon spectrometer, while stand alone muons are only reconstructed in the muon spectrometer. It makes it very simple to find the reconstruction efficiency of the inner detector:

1. The tag is found as a muon passing the usual requirements given in table 6.2.

2. The probe is found as a muon passing the same requirements except, that it is allowed to be both a combined muon or a stand alone muon. In addition it is not required to fire the trigger.

3. If the probe is combined, it is said to be passed.

The reconstruction efficiency of the inner detector is found as:

$$\epsilon_{\mu, \text{recon, ID}} = \frac{\text{Number of passed probes}}{\text{Number of all probes}}$$

The efficiency is shown in figure 6.18a and 6.18b as a function of $p_T$ and $\eta$ respectively. It gives an average efficiency of:

$$\epsilon_{\mu, \text{recon, ID}} = 0.992^{+0.00145}_{-0.00163}$$

This result is consistent with earlier analysis [65], which has found the efficiency of the inner detector to be $0.991 \pm 0.001$.

Measurement of the reconstruction efficiency using the three different chains and using the two periods separately, are found to be consistent. The plots are therefore not presented.
Efficiency of the muon spectrometer

Similar to the reconstruction efficiency of the inner detector, the efficiency of the muon spectrometer is found as the fraction of probes being reconstruction in the muon spectrometer with respect to all reconstructed muons. It is though not as simple, because muons are never reconstructed from the inner detector alone. Even segment tagged muons are required to have pixel hits in the muon spectrometer as described in section 5.1.6. The probe is therefore found among all tracks of the inner detector. To insure the probe to be the track of a real muon forming a Z boson with the tag, some very strong requirements need to be imposed. It includes requirements of pixel hits, SCT hits, hits in the B-layer, stronger requirements on the distance to the vertex and requirements on the significance of the reconstructed variables. In addition two cut are used on the seperation between the tag and the probe. All these requirements, which are inspired from earlier analysis [66][65] or recommended from the Heavy Ion Group at ATLAS, are given in table 6.8.

The efficiency of the muon spectrometer is found as described below:

1. The tag is found as a muon passing the usual requirements given in table 6.2.
2. The probe is found as a track passing the criterias given in table 6.8.
3. If the probe is matched to a muon using the requirements given in table 6.8, it is said to be passed.

The efficiency of the muon spectrometer is given as below:

\[ \epsilon_{\text{recon,MS}} = \frac{\text{Number of pass probes}}{\text{Number of all probes}} \]

The efficiency of the muon spectrometer is shown as a function of \( p_T \) and \( \eta \) in figure 6.19a and 6.19a. In average it gives an efficiency of:

\[ \epsilon_{\text{recon,MS}} = 0.966^{+0.00373}_{-0.00401} \]
Cuts

<table>
<thead>
<tr>
<th>Probe</th>
<th>Cuts</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>Number of pixel hits &gt; 1</td>
</tr>
<tr>
<td></td>
<td>Number of SCT hits &gt; 6</td>
</tr>
<tr>
<td></td>
<td>Number of B-layer hits &gt; 0</td>
</tr>
<tr>
<td></td>
<td>$d_0$, $z_0 \sin \theta &lt; 1.5$ mm</td>
</tr>
<tr>
<td></td>
<td>$</td>
</tr>
<tr>
<td></td>
<td>$</td>
</tr>
<tr>
<td></td>
<td>$\chi^2/n_{dof} &lt; 3$</td>
</tr>
<tr>
<td></td>
<td>$p^\text{cone20}_T/p_T^p &lt; 0.1$</td>
</tr>
<tr>
<td></td>
<td>pt $&gt; 10$ GeV</td>
</tr>
<tr>
<td></td>
<td>$\eta &lt; 2.4$</td>
</tr>
<tr>
<td></td>
<td>sign opposite to the tag</td>
</tr>
<tr>
<td></td>
<td>$</td>
</tr>
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<td></td>
<td>$</td>
</tr>
<tr>
<td></td>
<td>$\Delta\phi_{TP} &gt; 2$</td>
</tr>
<tr>
<td></td>
<td>$\Delta R_{TP} &lt; 4$</td>
</tr>
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</table>

Matching

<table>
<thead>
<tr>
<th></th>
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</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>same sign</td>
</tr>
<tr>
<td></td>
<td>$\Delta R_{TP} &lt; 0.2$</td>
</tr>
</tbody>
</table>

Table 6.8: Criterias for track selection in the inner detector and for matching tracks to muons.

Using the MUID muons and STACO muon, the reconstruction efficiencies of the muon spectrometer is found to be $0.967465 \pm 0.00363621$ and $0.958 \pm 0.001$, respectively. All chains are thereby shown to be consistent. It would be expected, though, if the Third Chain had a better reconstruction efficiency, since it is combined by the MUID and STACO algorithms. Earlier analysis [65] has found a result of $0.928 \pm 0.002$ for the first chain and $0.958 \pm 0.001$ for the second chain. All three chains is thereby found to have a slightly better reconstruction efficiency in this analysis.

(a) $p_T$ dependence.  
(b) $\eta$ dependence.

Figure 6.19: The single muon reconstruction efficiency in the muon spectrometer.
The total reconstruction efficiencies

When finding the efficiency of the inner detector and the muon spectrometer, the efficiency of combining the muons is already included. The total reconstruction efficiency is therefore given as the product of the efficiency of the inner detector and the efficiency of the muon spectrometer. The $p_T$ and $\eta$ dependence is given in figure 6.20a and 6.20b and the average value is found to be:

$$\epsilon_\mu^{\text{recon}} = 0.958^{+0.00509}_{-0.00555}$$

![Figure 6.20: The total single muon reconstruction efficiency.](image)

(a) $p_T$ dependence.  (b) $\eta$ dependence.

The reconstruction efficiency of the $Z$ boson is calculated as below:

$$\epsilon_Z^{\text{reco}} = 1 - (1 - \epsilon_\mu^{\text{reco}}(\eta_1))(1 - \epsilon_\mu^{\text{reco}}(\eta_2))$$

It is shown in figure 6.21 and the average value is found to be:

$$\epsilon_Z^{\text{reco}} = 0.998^{+0.0141}_{-0.0153}$$

6.7.3 Acceptance

The acceptance of the detector is found using $Z \to \mu \mu$ events generated by Pythia as described in section 6.1. Figure 6.22 shows the rapidity distributions of all generated $Z$ bosons formed by muons covering the full space. From this a rapidity plateau of $\sim 8$ is seen as expected from the calculation in section 3.4. This explain why no rapidity boost is seen in data. However, since the generated $Z \to \mu \mu$ events are produced with asymmetric beams, a rapidity boost would be expected in figure 6.22. Why this is not seen, is not well-known, probable it could be because the option of asymmetric beams has been ignored when generating the events.

The acceptance of the detector is found as the number of $Z$ bosons generated from muons, where at least one is within $|\eta| < 2.4$, divided by all $Z$ bosons. The
6.8 Cross Section

Having measured the efficiencies of the trigger system and the reconstruction algorithms, as well as the detector acceptance, it is possible to estimate the total number of $Z$ bosons and thereby to measure the differential $Z \rightarrow \mu\mu$ cross-section.
Since the efficiencies is not measured for each centrality group seperately, the corrections are only applied to the total number of measured Z bosons. To fulfill the centrality analysis, the efficiencies should also be found for each centrality group in the future.

The differential $Z \rightarrow \mu\mu$ cross-section $\sigma$ is found as:

$$\sigma = \frac{N_{\text{opposite-sign}} - N_{\text{same-sign}}}{A_Z \times C_Z \times \int L \, dt}$$

where $N_{\text{opposite-sign}}$ is the number of measured opposite-sign dimuons and $N_{\text{same-sign}}$ is the number of measured same-sign dimuons used as an estimate of the background. $A_Z$ is the acceptance, $C_Z$ is the combined efficiency of the trigger system and the reconstruction algorithms and $\int L \, dt$ is the integrated luminosity.

The differential $Z \rightarrow \mu\mu$ cross-section is thereby found to be:

$$\sigma = (1941 \pm 44.1) - (6 \pm 2.45)$$

where $N_{\text{opposite-sign}}$ is the number of measured opposite-sign dimuons and $N_{\text{same-sign}}$ is the number of measured same-sign dimuons used as an estimate of the background. $A_Z$ is the acceptance, $C_Z$ is the combined efficiency of the trigger system and the reconstruction algorithms and $\int L \, dt$ is the integrated luminosity.

The differential $Z \rightarrow \mu\mu$ cross-section is thereby found to be:

$$\sigma = \frac{(1941 \pm 44.1) - (6 \pm 2.45)}{0.450^{+0.00361}_{-0.00361} \times 0.998^{+0.0141}_{-0.0153} \times 0.926^{+0.0291}_{-0.0304} \times 25.5 \text{ nb}^{-1}} = 182^{+15.1}_{-13.4} \text{ nb}$$

Since the Monte Carlo sample for $Z \rightarrow \mu\mu$ production in p+Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV is not yet released, no direct comparison to theory is possible.

However, from the $Z \rightarrow \mu\mu$ cross-section measured in p+p collisions [56] and the relation between p+p and p+A cross-sections [67], it is possible to compare the obtained result.

The relation between a given cross-section measured in p+p and p+A collisions is given by [67]:

$$\sigma_{pA} = A R_{pA} \sigma_{pp} \quad \text{(6.1)}$$

where $\sigma_{pA}$ is the cross-section measured in p+A collisions, $A$ is the number of nucleons within the nuclei $A$, $R_{pA}$ is the nuclear modification factor of nuclei and $\sigma_{pp}$ is the cross-section measured in p+p collisions. This relation simply states, that the number of produced Z bosons in p+A collisions scales with the number of binary collisions, taken the nuclear modification factor into account. Assuming binary scaling, meaning that the nuclear modification factor is equal to one for Z production, the expression reduces to $\sigma_{pA} = A \sigma_{pp}$.

The cross-section depends strongly on the collision energy as seen in figure 6.24. This figure shows the differential cross-section of $Z \rightarrow \ell\ell$ for experiments operating with different collision energies. The experimental measurements are compared to a theoretical prediction based on NNLO QCD calculations using the MSTW 2008 PDF set as described in section 3.1. This prediction is shown for both p+p and p+\bar{p} collisions. For the ATLAS, CMS and CDF results, the differential cross-sections of the muon and electron channel are combined. It is also the case for the predictions. Since the brancing ratio between the muon and electron channel is almost equal - $0.2366 \pm 0.007\%$ and $3.363 \pm 0.004\%$, respectively [21] - the cross-section for the two channels are not distinguished.

By reading off the plot, the differential cross-section of $Z \rightarrow \mu\mu$ in p+p collisions at $\sqrt{s} = 5.02$ TeV is found to be about 0.6 nb. Combining this value with equation
Figure 6.24: The differential cross-section of $Z \rightarrow \bar{\ell}\ell$ as a function of collision energy measured in p+p collisions. [56]

6.1 using $A=207$ and assuming binary scaling, the measured cross-section is expected to be 124.2 nb.

The measured cross-section do thereby deviate from the expected value with about 50%. A final explanation for this huge deviation is not found, but small contributions are considered. Figure 6.25 shows the cross-section as a function of rapidity. In this plot, the measured cross-section is divided by $A=207$, and should therefore reflect the cross-section in p+p collisions assuming binary scaling. In addition a prediction of the differential $Z \rightarrow \mu\mu$ cross-section in p+p collisions at $\sqrt{s} = 5.02$ TeV is shown. The prediction is made from the PDF set, NNPDF2.1 NLO [68], which is a proton PDF made from NLO global fits as described in section 3.1. The rapidity distribution of the $Z$ bosons is computed using the MCFM code [69].

As seen from the rapidity distribution in figure 6.25, the deviation between the measured and predicted cross-section gets larger at low rapidity. Why it is so it not well understood.

If the $Z$ bosons produced at low rapidity were created by partons with a Bjorken-$x$ about 0.1, it could be explained by the shadowing effect. If it was the case, $Z$ bosons produced at high rapidity would be created by partons, where one would have a Bjorken-$x$ a bit smaller than 0.1 and the other one would have a Bjorken-$x$ a bit higher than 0.1. As seen in figure 4.11 the effect of antishadowing is peaking at $x=0.1$. $Z$ bosons produced at lower rapidity would therefore be stronger affected by the antishadowing effect compared to $Z$ bosons produced at high rapidity. In addition, since nuclear effects are stronger in central collisions compared to peripheral collisions, the effect of antishadowing could have explained why the number of measured $Z$ bosons per $N_{\text{coll}}$ produced in central collisions is relative larger than number of measured $Z$ bosons $N_{\text{coll}}$ produced in peripheral collisions - using the default Glauber Model. It was shown in figure 6.11.

However, this is not the case. As described in section 4.6, $Z$ bosons are pro-
CHAPTER 6. MEASUREMENT OF THE Z BOSON

Figure 6.25: Cross section as a function of rapidity both for data, which is divided by 207, and for a NNPDF2.1 prediction of the $Z \rightarrow \mu\mu$ cross section in p+p collisions at $\sqrt{s} = 5.02$ TeV. As seen a huge deviation exist and is seen to get larger at low rapidity.

duced from partons having a Bjorken-x of about 0.02 - an area which is not dominated neither by shadowing nor by antishadowing. The Z production is therefore expected to scale with the number of binary collisions and the rapidity distribution seen in figure 6.25 should therefore correspond to the prediction.

One thing, which could might contribute to the enhancement of the estimated number of Z bosons at low rapidity could be, if the trigger efficiency is underestimated. As seen in figure 6.17, the trigger efficiency is relative low at rapidities less than 0.7. However, the measured efficiency is consistent with earlier measurements and seem reliable, so the trigger efficiency could probably not explain the result either.

A small contribution to the enhancement of the measured cross-section could be explained by the fact, that coupling between Z bosons and down quarks are stronger than coupling between Z bosons and up quarks [21]. When comparing the measured cross-section with p+p results using binary scaling, the collisions involving neutrons are not taken into account. If only valence quarks are considered, the protons contain two up quarks and one down, while the neutrons contain two down quarks and one up. The axial and vectorial coupling of up quarks to Z bosons are $g_A^u = 0.50^{+0.04}_{-0.07}$ and $g_V^u = 0.29^{+0.10}_{-0.08}$, respectively, and the corresponding coupling of down quarks to Z bosons are $g_A^d = 0.524^{+0.050}_{-0.030}$ and $g_V^d = 0.33^{+0.05}_{-0.07}$. The coupling strength is proportional to $(g_A)^2 + (g_V)^2$ [12] which means, that the coupling of down quarks
to $Z$ bosons is enhanced by almost 15% compared to coupling of down quarks to $Z$ bosons. Since lead contains 82 protons and 125 neutrons in average, an enhancement of $Z$ production in p+Pb collisions of almost 10% would be expected compared to p+p collisions if the valence quarks was dominating.

However, at energies at the LHC, sea quarks are dominating. The protons and neutrons do therefore contain almost the same number of up and down quarks. The enhancement of $Z$ production in p+Pb collisions compared to p+p collisions caused by collisions involving neutrons, is therefore not expected to be significant.

Another reason could be found by considering double parton interactions. It has been shown in [70], that the number of double parton interactions in p+Pb collisions are increased by a factor of 600 compared to the p+p collisions. This is an effect, which is not included by equation 6.1. Measurement of $Z \rightarrow \mu\mu$ + 2 jets events at $\sqrt{s} = 7$ TeV has shown, that the fraction of double parton interactions relative to single parton interaction is \( f^{DP} = 0.02 \pm 0.014\) (stat.) \( \pm 0.009\) (sys.) [71]. Similar measurements has been done of $W \rightarrow l\nu + 2$ jet events [72]. The fraction of double parton interactions producing two $Z$ bosons is less likely, but since the fraction is increased by a factor of 600 in p+A collisions, it could maybe cause a small contribution. In this analysis, no events with two $Z$ bosons is observed. However, it could be because one of the two produced $Z$ bosons is decaying to muons with $\eta > 2.4$.

A part of the enhancement could also be caused by the measurement itself. The background is estimated to be 0.31% and do not include background sources, which produce same-sign dimuons. From [56] the background from sources producing same-sign dimuons was estimated to be almost in the same order as the QCD background as described in section 6.3. However, it will only decrease the measured cross-section by around 0.31%, and is therefore not significant.

In addition the efficiencies could be underestimated of some reason, but whether this is the case or not is hard to say. Compared to earlier measurements, the efficiency do not seem to be underestimated.

Summing up; the measured differential $Z \rightarrow \mu\mu$ cross-section is found to be 182 nb. The expected cross-section obtained from p+p collisions assuming binary scaling is found to be 124 nb. Contributions to this deviation includes the effect from neutron collisions and double parton interactions. However, these effects can only explain a very small part of the deviation. The enhancement of $Z$ production in p+A collisions compared to p+p collisions assuming binary scaling can therefore not be explained without further investigation in the future. Maybe it could be explain by nuclear effect or other unforeseen influences causing deviation from binary scaling.
Chapter 7

Conclusion

This thesis has described the measurement of $Z$ bosons produced in p+Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV using the ATLAS detector at the LHC. p+Pb collisions are performed with asymmetric beams causing a rapidity boost of 0.47. This boost was not observed in data for unknown reasons. The $Z$ bosons - produced by the Drell-Yan process - was measured in the muonic decay channel. From an integrated luminosity of 25.5 nb$^{-1}$, 1941 $Z$ bosons was measured and the background, obtained from the number of same-sign dimuons, was found to be 0.3%.

The measurement was performed within different centrality groups, to measure the centrality distribution. From this distribution, the $Z$ production has been found not to scale with the number of binary collisions using the default Glauber Model. It is in contrast to an earlier measurement of $Z$ bosons in Pb+Pb collisions [39]. Instead it was shown, that a correct description of the centrality in p+Pb collisions might require the Glauber-Gribov Model, which includes fluctuations in the nucleon-nucleon cross-section.

In addition to the centrality analysis, the differential $Z \rightarrow \mu\mu$ cross-section has been measured. To estimate the total number of produced $Z$ bosons, the efficiencies of the trigger system and the reconstruction algorithms, as well as the detector acceptance was found. The efficiencies of the trigger system and the reconstruction algorithms were performed datadriven by investigating single muons and by using the tag and probe method, respetitely. The acceptance of the detector was found with a Monte Carlo sample generated by Pythia 8. The results was found to be:

$$\epsilon_{\text{trig}}^Z = 0.926^{+0.0291}_{-0.0304}$$
$$\epsilon_{\text{reco}}^Z = 0.998^{+0.0141}_{-0.0153}$$
$$A_Z = 0.450^{+0.00361}_{-0.00361}$$

Using these efficiencies and the acceptance, the differential $Z \rightarrow \mu\mu$ cross-section was found to be:

$$\sigma = 182^{+15.1}_{-13.4} \text{ nb}$$

No earlier measurements of $Z$ production in p+A collisions exist and in addition, the Monte Carlo simulation of $Z \rightarrow \mu\mu$ in p+Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV is not yet
released. Comparison of the cross-section has therefore only been performed with measurements in p+p collisions. By the assumption of binary scaling and absence of nuclear effects, the relation between the cross-section in p+p collisions and p+Pb collisions is $\sigma_{pPb} = 207\sigma_{pp}$. Using this relation, the measured cross-section was found to deviates with about 50%. A final explanation for this large deviation has not been found. A small contribution could be caused by the stronger coupling of Z bosons to down quarks than of Z bosons to up quarks, which could have an effect since p+Pb collisions involve collisions with neutrons in contrast to p+p collisions. In addition, the number of double parton interactions is increased by a factor of 600 in p+A collisions compared to p+p collisions. However, these contributions are found not to be significant.

Another contribution to the large cross-section could be caused by nuclear effects - more specific, the effect of antishadowing, which could cause an increase of Z produced especially in central collisions. From nuclear PDFs and earlier measurement, this is not expected and will need further investigation.

So, the most interesting results obtained from this thesis is not concluding, but gives new related questions.

Is the production of Z bosons affected by nuclear effects or not? and is centrality measurements in p+A collisions described by the default Glauber Model or is the Glauber-Gribov Model, which includes fluctuations in the nucleon-nucleon cross-section, required?

To get the answer to these quetions, further investigation must be performed.
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