

SEARCH FOR MAGNETIC MONOPOLES AND STABLE
HIGH-ELECTRIC-CHARGE OBJECTS IN 13 TEV CENTRE-OF-MASS
ENERGY PROTON-PROTON COLLISIONS WITH THE ATLAS
DETECTOR AT THE LHC

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Abstract

One asymmetry of the current formulation of classical electrodynamics is the absence of a magnetic monopole, despite the existence of isolated electric charges. The existence of such a particle with net magnetic charge not only restores the symmetry between electricity and magnetism but also provides an explanation for the quantization of electric charges due to a correlation between the magnitude of the fundamental magnetic charge and that of the fundamental electric charge. Magnetic monopoles do not decay and lose kinetic energy in matter similarly as ions with high electric charge. This work presents the strategy and results of a search for magnetic monopoles and stable high-electric-charge objects with mass up to 4000 GeV. It is performed with 138 fb^{-1} of proton-proton collision data at a centre-of-mass energy of $\sqrt{s} = 13 \text{ TeV}$, collected during 2015 – 2018 with the ATLAS Detector at the Large Hadron Collider. No deviation from Standard Model predictions is observed, thus exclusion limits are placed on production models. These limits surpass the ones obtained in the previous search. With the increase in statistics and an interest in heavy-ion collisions, the search for magnetic monopoles and stable high-electric-charge objects will continue to play an important role in the quest for physics beyond the Standard Model.

To my parents, who support me for every decision I make. It is to them that this thesis is dedicated.

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Abbreviations

AOD Analysis Object Data

ASDBLR Amplifier-Shaper-Discriminator Baseline Restorer

ATLAS A Toroidal LHC ApparatuS

CL Confidence Level

CLs Method to estimate modified frequentist confidence limits

CSCs Cathode Strip Chambers

DAOD Derived Analysis Object Data

DESD Derived Event Summary Data

DESDM_EXOTHIP Derived Event Summary Data for HIP analysis

DTMROC Drift-Time Measuring Readout Chip

DY Drell-Yan

GRL Good Runs List

EM Electromagnetic

EM1 The first layer of the EM Calorimeter

EM2 The second layer of the EM Calorimeter

EM3 The third layer of the EM Calorimeter

ESD Event Summary Data

FCAL Forward Calorimeter

GUT Grand Unified Theories

HCAL Hadronic calorimeter
HECO High-electric-charge object
HIP Highly ionizing particle
HITS Particle energy deposits in simulation
HLT High-Level Trigger
HT High threshold
IBL Insertable B-Layer
ID Inner Detector
IP Interaction point
L1 First-Level trigger
LAr Liquid argon
LHC Large Hadron Collider
LT Low threshold
MC Monte Carlo (simulation)
MDTs Monitored Drift Tubes
MS Muon Spectrometer
PAI Photo-Absorption Ionization model
PDF Parton distribution function
PF Photon fusion
QCD Quantum Chromodynamics
QED Quantum Electrodynamics
RAW Raw data format
RDO Raw Data Output
ROD Readout driver, i.e., backend electronics
RoI Region of interest, i.e., a subregion of the ATLAS Detector used for HLT
RPCs Resistive Plate Chambers

SCT SemiConductor Tracker

SM Standard Model

T0 Tier 0

T1 Tier 1

T2 Tier 2

TCGs Thin Gap Chambers

TDAQ Trigger and Data Acquisition

TR Transition Radiation

TRT Transition Radiation Tracker

UPC Ultraperipheral collisions

WLCG Worldwide LHC Computing Grid

Chapter 1

Introduction

Magnetic monopoles and stable high-electric-charge objects, collectively known as highly ionizing particles (HIPs), have long been pursued in Beyond the Standard Model physics research. In order to motivate and understand the work described in this thesis, this chapter begins with an exposition of theories or models that predict HIPs. It goes on to review searches at colliders and astrophysical experiments.

1.1 Magnetic Monopoles

There has been a long-running dispute about the possibility of such an object known as a magnetic monopole. With distributions of electric volume charge densities ρ_E and current densities $\mathbf{j}_E$, the time evolution of electric fields $\mathbf{E}$ and magnetic fields $\mathbf{B}$ is described in Maxwell's Equations:

$$\begin{aligned}\nabla \cdot \mathbf{E} &= \rho_E \\ \nabla \cdot \mathbf{B} &= 0 \\ \nabla \times \mathbf{E} &= -\frac{\partial}{\partial t} \mathbf{B} \\ \nabla \times \mathbf{B} &= \frac{\partial}{\partial t} \mathbf{E} + \mathbf{j}_E.\end{aligned}\tag{1.1}$$

Although Equation 1.1 can be readily modified to include magnetic sources, the existence of a magnetic monopole invites justification. On the one hand, this possibility is favored because Maxwell's Equations do not remain invariant under duality transformations given by $\theta \in [0, 2\pi]$ that transform the fields $\mathbf{E}$ and $\mathbf{B}$ as follows:

$$\mathbf{E} - i\mathbf{B} \rightarrow e^{i\theta}(\mathbf{E} - i\mathbf{B}). \quad (1.2)$$

This is ultimately a consequence of the absence of magnetic monopoles, which breaks the electric-magnetic dual symmetry Maxwell's Equations should otherwise possess. On the other hand, the construction of a magnetic monopole was not trivial, as its absence, i.e., $\nabla \cdot \mathbf{B} = 0$, was assumed a necessity in order to define the magnetic field $\mathbf{B} = \nabla \times \mathbf{A}$ with the vector potential $\mathbf{A}$, such that

$$\nabla \cdot (\nabla \times \mathbf{A}) = 0. \quad (1.3)$$

Very few believed in the existence of a magnetic monopole [1] until the birth of quantum mechanics. Since then, this possibility has been revisited over and over again [2–17]. As reviewed in [18–23], three key monopole types bear emphasis. The first type was put forward by Dirac [2, 3], who proved the consistency of magnetic monopoles between quantum mechanics and the theory of electromagnetism. The next type came when 't Hooft and Polyakov independently postulated a more generalized form of magnetic monopole formulated in quantum field theory [4, 5]. The 't Hooft-Polyakov monopole exists in all Grand Unified Theories and is fundamentally different from the Dirac monopole. The third monopole was developed by Cho and Maison in the electroweak sector of the Standard Model of Particle Physics as a generalization of the Dirac monopole [7, 15]. In order to gain insight into these magnetic monopoles, I will first introduce the concept of gauge symmetry.

1.1.1 Gauge Symmetry

Many successful theories of physics are built on the foundation of gauge symmetry. One example is classical electrodynamics, where magnetic fields are explicitly defined by vector potentials. When varied in space, a magnetic vector potential gives rise to both electric and magnetic fields. There are infinitely many vector potentials that give rise to the same electric and magnetic fields. These potentials are referred to as gauge fields, denoted by $\mathbf{A}$, and are related by gauge transformations χ :

$$\mathbf{A} \rightarrow \mathbf{A} - \nabla\chi, \tag{1.4}$$

where the gauge function χ satisfies $\nabla^2\chi = 0$. Since the resulting physical observables, i.e., electric and magnetic fields, do not change under the gauge transformations, classical electrodynamics is said to have a gauge symmetry.

Despite seeming to be a mathematical construct (see Equation 1.3) for the definition of magnetic fields, vector potentials possess physical meaning that becomes manifest in quantum mechanics, where a particle is described by its wave function, denoted as ψ . To couple electromagnetism to quantum mechanics, one replaces the momentum operator $\mathbf{p}$ by the canonical momentum operator with the minimal coupling between the electric charge and the (time-independent) vector potential:

$$\mathbf{p} = -i\hbar\nabla \rightarrow -i(\hbar\nabla - i\frac{e}{c}\mathbf{A}),$$

so that the Schrodinger equation is recast as

$$i\frac{\partial\psi}{\partial t} = [-\frac{1}{2m}(\hbar\nabla - i\frac{e}{c}\mathbf{A})^2]\psi.$$

The Schrodinger equation is invariant under the gauge transformations in Equation 1.4 if and only if the wave function ψ changes by a phase $\exp(-i\frac{e}{\hbar c}\chi)$:

$$\psi \rightarrow \exp\left(-i\frac{e}{\hbar c}\chi\right)\psi, \quad (1.5)$$

which leaves invariant the probability density $|\psi|^2$ of the wave function ψ . Hence, for the introduction of magnetic source charges to be compatible with quantum mechanics, the vector potential due to the magnetic charge must be a gauge transformation, i.e., $-\nabla\chi$. Since the resulting phase can be removed upon a gauge transformation, there is no change of physical fields, thus the monopole stays undetected.

The above gauge symmetry is obtained from a $U(1)$ gauge transformation as per Equation 1.5. I will show explicitly that the $U(1)$ gauge symmetry underlies the existence of the Dirac monopole. This symmetry is also important in the construction of the 't Hooft-Polyakov monopole and the Cho-Maison monopole. It is an interesting subject on its own as to why particles transform the way they do, but this will not be discussed in this work, which will focus on the consequences of these gauge transformations.

1.1.2 Dirac Monopole

It is important to note that although the vector potential, also known as a gauge field, is often assumed regular everywhere, this is not a mandatory requirement, as the gauge field is not a physical observable that can be measured. Having realized this, Dirac argued the potential existence of a magnetic monopole of charge g associated with a magnetic field $\mathbf{B}$, such that

$$\nabla \cdot \mathbf{B} = 4\pi g\delta(\mathbf{r}), \quad (1.6)$$

which reduces to the expression in Equation 1.3 for all positions $\mathbf{r}$ except for where the monopole is located due to the Dirac delta function $\delta(\mathbf{r})$. Furthermore, Dirac was able to justify how the monopole stays undetected and consequently draw implications about this feature.

In order to detect a magnetic monopole, an experiment one can devise is to keep the monopole fixed in space and move an electron along an arbitrary closed path Γ that goes around the monopole. The monopole's presence in the form of a gauge potential $\mathbf{A}$ modifies the action of the electron. That is, the monopole is undetected if the action is invariant. Given the Lagrangian of the electron with mass m and charge e

$$\mathcal{L} = \frac{1}{2}m\dot{\mathbf{r}}^2 + \frac{e}{c}\dot{\mathbf{r}} \cdot \mathbf{A},$$

where c denotes the speed of light and the vector potential is a gauge transformation, i.e., $\mathbf{A} = -\nabla\chi$, the action reads

$$S = \int dt \mathcal{L} = \frac{1}{2} \int dt m \dot{\mathbf{r}}^2 + \frac{e}{c} \oint_{\Gamma} d\mathbf{r} \cdot \mathbf{A} = S_0 + \delta S.$$

Clearly, this action, in particular the term δS , is not generally invariant for an arbitrary gauge function χ because it gives rise to a nontrivial phase

$$\epsilon = \exp\left(\frac{i}{\hbar c} \delta S\right) = \exp\left(\frac{i}{\hbar c} e \int_{\Gamma} d\mathbf{r} \cdot \mathbf{A}\right)$$

in the probability amplitude for the electron to travel from one spacetime point to another. However, under one condition the magnetic monopole will stay undetected. That is when the change in phase is

$$\frac{e}{\hbar c} \oint_{\Gamma} d\mathbf{r} \cdot \mathbf{A} = \frac{e}{\hbar c} \int d\mathbf{S} \cdot \mathbf{B} = \frac{4\pi e g}{\hbar c}, \quad (1.7)$$

where

$$\frac{4\pi eg}{\hbar c} = 2\pi n \quad \text{for an arbitrary integer } n, \quad (1.8)$$

is known as the Dirac quantization condition [2]. A rearrangement of Equation 1.8 yields

$$eg = \frac{n}{2}\hbar c. \quad (1.9)$$

Taking $n = 1$ and dividing both sides of Equation 1.9 by the fine structure constant $\alpha \approx \frac{1}{137}$, one obtains the ratio between a unit magnetic charge g_D and the fundamental electric charge e :

$$\frac{g_D}{e} = \frac{1}{2\alpha} \approx 68.5. \quad (1.10)$$

Another interesting feature of the Dirac monopole is the vector potential, known as the Dirac potential. Its uniqueness has been mentioned earlier as not being regular everywhere. Shown in Figure 1.1, the Dirac potential is singular along the line known as the Dirac string that extends from (yet excluding) the monopole to infinity, leading to a singular flux of magnetic field $\mathbf{B}_{\text{sing}}$. Dirac held that a point-like magnetic monopole is allowed if it is attached to the end of the Dirac string and satisfies the Dirac quantization condition per Equation 1.9. However, the Dirac string poses a challenge when attempting to promote the Dirac monopole from quantum mechanics to quantum field theory. The string singularities are not covariant and thus violate the principle of quantum field theory that is the expectation that physics should be independent of the choice of spacetime coordinates.

Nevertheless, the Dirac monopole leaves two main legacies. First, by introducing magnetic volume charge densities ρ_M and current densities $\mathbf{j}_M$ that transform with their electric counterparts as

$$\begin{aligned} \rho_E - i\rho_M &\rightarrow e^{i\theta}(\rho_E - i\rho_M), \\ \mathbf{j}_E - i\mathbf{j}_M &\rightarrow e^{i\theta}(\mathbf{j}_E - i\mathbf{j}_M), \end{aligned} \quad (1.11)$$

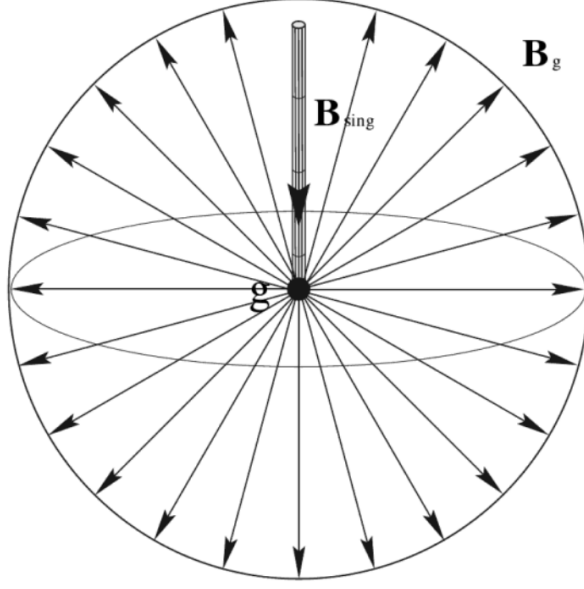


Figure 1.1: Depiction of the Dirac string and its associated magnetic field $\mathbf{B}_{\text{sing}}$ [19].

the Dirac monopole restores the electric-magnetic dual symmetry of Maxwell's Equations with

$$\begin{aligned}
 \nabla \cdot \mathbf{E} &= \rho_E \\
 \nabla \cdot \mathbf{B} &= \rho_M \\
 \nabla \times \mathbf{E} &= -\frac{\partial}{\partial t} \mathbf{B} - \mathbf{j}_M \\
 \nabla \times \mathbf{B} &= \frac{\partial}{\partial t} \mathbf{E} + \mathbf{j}_E
 \end{aligned} \tag{1.12}$$

left invariant under the transformations in Equation 1.2. The second legacy is that the Dirac monopole implies the quantization of magnetic charges, as a Dirac monopole would carry a magnetic charge g that is a multiple of the unit charge g_D . Also, Equation 1.10 implies an approximate equivalence between a monopole with one Dirac unit of magnetic charge and an ion with electric charge $|z| = 68.5$. The consequence of this equivalence is of critical importance and merits further examination in later chapters.

1.1.3 't Hooft-Polyakov Monopole

The list of challenges in treating the Dirac monopole in quantum theories continues. Its mass and spin are completely unconstrained. In addition, it cannot be handled with perturbation theory, since the Dirac quantization condition in Equation 1.8 implies that electric and magnetic couplings are inversely correlated, hence, cannot be simultaneously small. These problems disappear if monopoles are not fundamental particles but field configurations with finite size and energy. Initially, it was proven that many gauge theories contain such monopoles [4, 5]. Later, it was discovered that Grand Unified Theories (GUT), originally proposed to unify the Standard Model interactions, in general predict these monopole solutions [24]. These solutions manifest as stable knots in the Higgs fields with a net magnetic charge. They necessarily exist in Grand Unified Theories because they are needed to make the energy of the GUT Higgs-field configuration finite.

To illustrate the inevitability of such monopoles, consider a toy model with a real scalar Higgs field, Φ^a ($a = 1, 2, 3$), whose potential energy density has a non-zero minimum v . Any Higgs state where the potential energy density is minimized is known as a vacuum. There are infinitely many vacua; each is connected to another via a gauge transformation of the Higgs field. The Higgs field has interesting properties and implications if it is assumed to be finite in energy. First, finite energy dictates that the field's potential energy density must be minimized at spatial infinity. That is, any Higgs field configuration is a map that takes the sphere at spatial infinity to a sphere of its vacuum states. In the case where the field orientation depends on spatial coordinates ($\vec{r}$), the field can take the form

$$\Psi^a(\vec{r}) = v f(r) \hat{r}^a, \quad (1.13)$$

where $r = |\vec{r}|$, $\hat{r}^a$ is one component of the unit vector $\hat{r} = \frac{\vec{r}}{r}$, and $f(r)$ is a function that is subject to two boundary conditions: $f(0) = 0$ and $f(\infty) = 1$. Sketched in Figure 1.2, this

configuration is known as the *hedgehog configuration* [5], where the Higgs field is radially directed outwards from the origin. Found in a vacuum at large distances, this Higgs-field configuration is proven to be stable if the boundary condition at infinity is fixed and if the field is continuous [25]. It results in a localized energy around a point that behaves as a particle. Furthermore, this particle is magnetically charged, following from the finite-energy

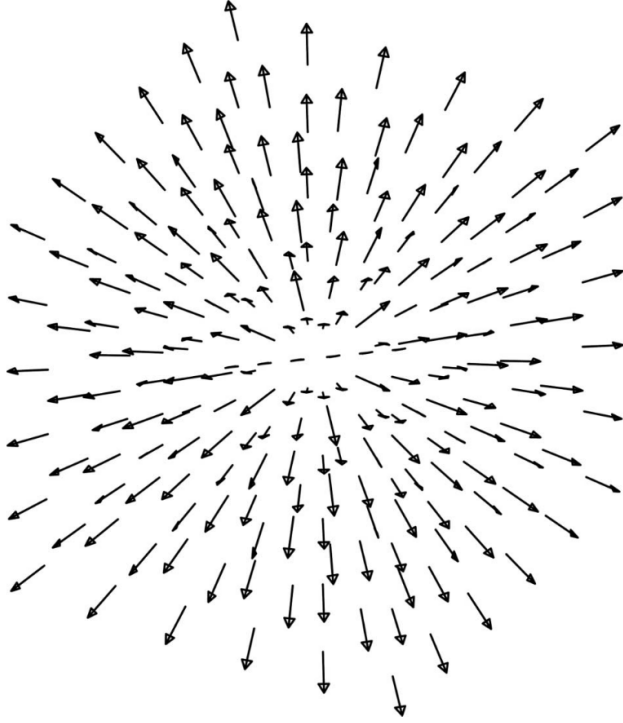


Figure 1.2: Depiction of the Higgs field Φ^a in the hedgehog configuration [22]. Each arrow represents the Higgs field at a given point in space.

assumption of the Higgs field. It can be shown that the total energy of the Higgs field in Equation 1.13 has a linear dependence on r , which in turn leaves the total energy infinite. If gauge fields of a magnetic charge are introduced as per Equation 1.6, then the magnetic charge contributes to the dynamics at the spatial origin yet does not alter the Higgs vacuum at the spatial infinity. In addition, the magnetic charge is quantized with its magnitude fixed by the Higgs field solution considered. In the case of Equation 1.13, this charge equals to

one Dirac charge (g_D).

Unlike the Dirac monopole whose mass is completely unknown, the 't Hooft-Polyakov monopole in Grand Unified Theories has a mass on the order of 10^{18} GeV, about two orders of magnitude higher than the GUT energy scale. However, the energy scale can be significantly lowered, for example through extra dimensions, giving rise to monopole masses possibly accessible at the Large Hadron Collider (LHC) [6].

1.1.4 Cho-Maison Monopole

Unlike Grand Unified Theories, the Standard Model (SM) does not contain the 't Hooft-Polyakov monopole since the SM Higgs field, a complex scalar field, does not support the spherically symmetric hedgehog monopole configuration (see Figure 1.2). Nevertheless, the Standard Model predicts its own monopole that arises from the SM Higgs field within the electroweak sector, according to Cho and Maison [7]. This monopole, in essence, is an alternative scenario of the 't Hooft-Polyakov monopole. It is also constructed as a localized energy associated with the Higgs field, by viewing the SM Higgs field as a map illustrated in Figure 1.3 from a Riemann sphere, i.e., an extended complex plane formed by a complex plane and a point at infinity. The localization of energy behaves as a particle with a net magnetic charge. The mass of the Cho-Maison monopole, computed as the total energy of the SM Higgs field configuration, has a term that diverges at the origin. Therefore, the monopole mass in the Standard Model is infinite. However, extensions of the Standard Model have been considered to regulate the monopole mass and constrain it within a range accessible to colliders [8, 10–12].

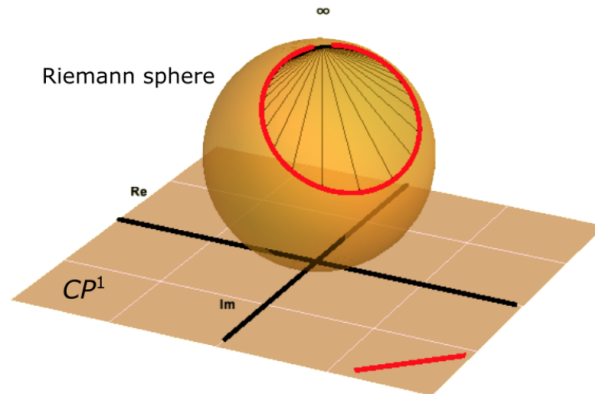


Figure 1.3: Depiction of the SM Higgs field as a complex projective line from the Riemann sphere [23].

1.2 High-Electric-Charge Objects

Apart from magnetic monopoles, high-electric-charge objects (HECOs) are also strongly motivated [26–30]. As the name suggests, HECOs carry electric charges greater than that of all electrically charged fundamental particles known to date. They remain stable and do not decay over time into objects with smaller charge because the high-charge configuration is favored under certain circumstances.

1.2.1 Q-Balls

One example of HECOs is Q-balls [27], which are stable and localized field configurations. As aggregates of a number of particles, Q-balls are stable, i.e., do not decay, because their energy per unit charge is smaller than any other configuration.

To see how Q-balls arise, consider a theory of two real scalar fields, denoted as ϕ_1 and ϕ_2 , respectively, whose potential energy density $V(\phi)$ is a function of $\phi = \sqrt{\phi_1^2 + \phi_2^2}$. If the fields are assumed to be symmetric under $SO(2)$ transformation, they will find solutions of

the following form

$$\begin{aligned}\phi_1 &= \phi(r) \cos \omega t, \\ \phi_2 &= \phi(r) \sin \omega t,\end{aligned}\tag{1.14}$$

where ω is the angular velocity of a field rotation in the $SO(2)$ space. Assuming $\phi(r)$ is constant in some volume of radius R and zero elsewhere, the conserved charge reads

$$Q = \int d^3\mathbf{x} (\phi_1 \partial_0 \phi_2 - \phi_2 \partial_0 \phi_1) = \frac{4\pi}{3} R^3 \omega \phi^2.\tag{1.15}$$

To prove that these two scalar fields can form a Q-ball, it suffices to find a solution of Equation 1.14 that minimizes the energy

$$E = \frac{4\pi}{3} R^3 \left(\frac{1}{2} \omega^2 \phi^2 + V \right)\tag{1.16}$$

and conserves the $SO(2)$ charge in Equation 1.15. First, the energy in Equation 1.16 takes a minimum value of $E_0 = Q \sqrt{2V/\phi^2}$. Moreover, the Q-ball will not decay in order to conserve charge, if the Q-ball's energy per unit charge, i.e., $E_0/Q = \sqrt{2V/\phi^2}$, is smaller than the mass of the scalar particle.

Q-balls can be realized in supersymmetric generalizations of the Standard Model [31], as they have scalar fields to provide conserved charges. A supersymmetric Q-ball can either be electrically charged [32] or neutral [33], but it has to be electrically charged to be a HECO.

1.2.2 Quark Matter

Quark matter is a hypothetical form of baryonic matter (matter with net baryon number) composed entirely of quarks. It is predicted to be more stable than nuclei when certain conditions are met. Its stability follows from energy considerations such that quark matter, if it exists, could be the true ground state of baryonic matter.

The first quark matter considered was strange quark matter, where the addition of strange quarks to up- and down-quark-containing hadrons was investigated [26, 34]. It was argued that the strange quark matter should be stable if the three quark flavors are approximately equal in number in the hadron, despite the fact that all known particles with strange quarks are unstable, as the strange quark decays to an up quark via the weak interaction. The argument is predicated on the Pauli Exclusion Principle, which implies that strange quarks will occupy their lowest possible energy states that are more energetically favorable than the unoccupied energy states of up and down quarks. The larger the number of strange quarks, the more suppressed the conversion of a strange quark to an up quark from the energy barrier the strange quark has to overcome to go to a higher energy state. Therefore, such a system where strange quarks coexist with up and down quarks has a lower chemical potential, i.e., a change of energy from the conversion of strange quarks to up quarks, than one with only up and down quarks. The stability of strange quark matter was predicted for a wide range of baryonic numbers (A) from a few to 10^{57} [35], where $A = (n_u + n_d + n_s)/3$, with n_u , n_d and n_s being the number of up, down and strange quarks in the strange quark matter. A more recent study explored the potential of quark matter with only up and down quarks, known as the up-down quark matter [30]. With an improved modelling of the quark flavors, it was shown that up-down quark matter is more stable than strange quark matter. Furthermore, the up-down quark matter is more stable than nuclei for baryon number $A = (n_u + n_d)/3 \geq 300$, which protects them from decaying into ordinary nuclei (with $A < 300$). The large positive baryon number also leads to a large positive electric charge of $Z \geq 100$, as shown in Figure 1.4.

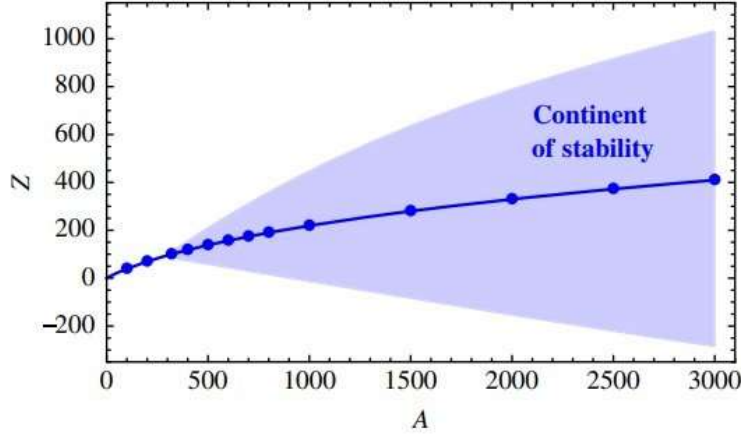


Figure 1.4: The electric charge (Z) of up-down quark matter as a function of baryon number (A) [30].

1.3 Previous Experimental Searches

As reviewed in the preceding section, magnetic monopoles and HECOs find strong motivation in a variety of theories. Consequently, enthusiasm for magnetic monopole and HECO searches emerged in recent decades in collider experiments [20, 23, 36, 37] and in cosmic rays [23, 37–39]. According to Equation 1.10, a magnetic monopole of one fundamental magnetic charge g_D [3] loses energy in matter via ionization similarly to an ion with electric charge $|z| = 68.5$ [40–42]. Therefore, magnetic monopoles and HECOs are referred to as highly ionizing particles (HIPs). Unlike particular names, which evoke distinct theoretical underpinnings of these particles, the collective term “HIP” emphasizes the highly ionizing nature they share. To maximize sensitivity to these particles, experimental searches target a wide HIP kinematic regime, by adopting one of the following approaches when it comes to the HIP production models.

The first approach makes no assumption on the HIP production model and thus no prediction of their theoretical production cross sections. The upper limits of these cross sections

can be derived from inclusive production in fiducial regions, where particle kinematics are restricted to yield high and uniform detector acceptance and reconstruction efficiency. These searches are naturally disposed to perform a brute-force scan of the HIP parameter space in a highly model-independent way. Furthermore, they tend to set stringent bounds on the cross sections, as a result of defining fiducial regions. The other two approaches entail minimal models with the fewest possible parameters. On the one hand, these production models prove useful in the face of non-trivial detector acceptance that impedes the construction of fiducial regions. They serve to predict the HIP kinematic distributions, which are used to model detector response through simulation. The simulated detector response of the HIPs is then tested against a set of signal selection criteria. The resulting signal selection efficiency translates to an upper limit for the measured cross section. On the other hand, these models are necessary to obtain mass limits, which early searches did not take an interest in, due to limited experimental reach to a wide range of masses. Assuming at least one model, the theoretical cross sections can be computed for HIPs with different masses, charges, and perhaps other quantum numbers. Subsequently, the HIPs are excluded if they have masses less than a mass limit, beyond which the theoretical cross section drops below the cross-section limit such that the experiment is no longer sensitive to the presence of HIPs.

1.3.1 Collider Searches

Different collision systems, such as electron-positron collisions, lepton-hadron collisions, hadron-hadron collisions and heavy-ion collisions, have been employed by the numerous collider experiments for HIP searches. What distinguishes a collision system is its effective collision energy and what production mechanisms it can possibly support, both of which strongly impact the HIP kinematic properties. To date, stable Dirac magnetic monopoles have been pursued in all four collision systems, while the hunt for HECOs is exclusive to

hadron-hadron and heavy-ion collisions.

Searches in Electron-Positron Collisions

The searches for magnetic monopoles in electron-positron collisions were beginning to take shape in the 1980s, assuming pair production from the annihilation of the collision particles. Hindered by the available collision energies on the order of 100 GeV, these searches were preoccupied with setting cross-section limits instead of mass limits. For example, early searches by experiments at DESY's PETRA accelerator [43] and SLAC's PEP accelerator [44, 45] and the CLEO Detector at Cornell University [46] were not sensitive to monopoles with masses larger than 13 GeV. The mass range was later extended by the TASSO Detecor at PETRA [47] and the TRISTAN Ring at KEK [48, 49], yet the collision energies below 100 GeV remained a main limitation until the Large Electron-Positron Collider (LEP) at CERN forged ahead with its search program in the 1990s. The MODAL Experiment at LEP set its sights on lower charges from $0.2g_D$ to $1g_D$ [50], and later the L6-MODAL Experiment at LEP-1 was sensitive to monopoles with magnetic charges of $1g_D$ and $2g_D$ [51], despite the fact that neither was able to probe monopoles with masses greater than 45 GeV. The strongest exclusion limit was set by the OPAL Experiment at LEP-2, which searched for $1g_D$ monopoles in a 62.7 pb^{-1} collision dataset at a centre-of-mass energy of 206 GeV on average [52]. The result was interpreted as an upper cross-section limit of 0.05 pb for masses from 45 to 102 GeV.

Searches in Lepton-Hadron Collisions

With the electron-positron collision programs being a resounding success, a unique asymmetric collision system came about in the early 1990s at the HERA electron (positron)-proton collider. Shortly after, the H1 Experiment performed a search for magnetic monopoles, predicated on the assumption of pair production via photon fusion, at a centre-of-mass energy of

300 GeV [53]. According to the production model, spin-0 monopoles were produced by elastic interactions between two photons that were radiated from the positron and the proton, while in contrast spin- $\frac{1}{2}$ monopoles were produced by inelastic collisions between a photon from the positron and one from a quark in the proton. The resulting cross-section upper limits varied between 0.05 pb and 2 pb for monopoles of charges up to $6g_D$ and masses up to 140 GeV.

Searches in Hadron-Hadron Collisions

In comparison to the two collision systems mentioned above, hadron colliders have a longer life span and have undoubtedly become the energy frontier ever since the startup of the CERN Intersecting Storage Rings (ISR) in the 1970s. While the monopole searches were, to a large extent, carried out in proton-proton (-antiproton) collisions, the HECO searches were restricted to heavy-ion collisions until they were performed in the LHC proton-proton collisions.

There was an explosion of interest for the searches of magnetic monopoles at the ISR, the Super Proton Antiproton Synchrotron, and the Tevatron proton-antiproton collider in the past decades [54–65]. Among these searches, the Tevatron searches stand out as particularly noteworthy, as they had provided the most comprehensive limits after it began operating in the 1980s. Performed at a centre-of-mass energy of 1.8 TeV, early searches for monopoles trapped in matter surrounding the D0 Detector at Tevatron [55, 56] carried great weight, because they opened the door for probing monopoles in hadronic collisions and for the first time interpreted the results by the Drell-Yan mechanism [66], i.e., quark-anti-quark annihilation to a virtual photon that decays to a pair of monopoles, which became a main benchmark model for further searches in hadron colliders. The upper limits on the cross sections were found to be 0.6 pb, 0.2 pb, 0.07 pb and 0.2 pb for $1g_D$, $2g_D$, $3g_D$ and $6g_D$ monopoles, respectively, whose corresponding mass limits were set to 265 GeV, 355 GeV,

410 GeV and 375 GeV [55, 56]. Later in 2006, the CDF Experiment at Tevatron performed another search with collisions at a higher centre-of-mass energy of 1.96 TeV, where an upper cross-section limit of 0.2 pb was attained for monopoles with masses between 200 GeV and 700 GeV, followed by a lower mass limit of 360 GeV [54].

The twenty-first century is notable for the commissioning of the LHC to operate at unprecedented collision energies, such that the last decade witnessed considerable ingenuity and efforts made for monopole searches. The first search was carried out by the ATLAS Collaboration in 2012 by virtue of 7 TeV proton-proton collision data collected in Run 1 [67] and that served as a prelude to later LHC searches of not only monopoles but HIPs in general [68–83]. The most stringent bounds were placed by the ATLAS [67–74, 84] and the MoEDAL Experiments [77–83] from the outset. In ATLAS, the most recent analysis drew on 34.4 fb^{-1} 13 TeV proton-proton collision data collected during early Run 2 (2015 – 2016) [72]. Assuming the Drell-Yan mechanism, this analysis set mass limits of 1725 – 2370 GeV for spin-0 and spin- $\frac{1}{2}$ monopoles with magnetic charges of $1 - 2g_D$ and mass limits of 1355 – 2050 GeV for spin-0 and spin- $\frac{1}{2}$ HECOs with charges up to $|z| = 100$. As for the MoEDAL Experiment [80, 82, 83], it analyzed 4.0 fb^{-1} of proton-proton collision data recorded at the same centre-of-mass energy, probing an additional spin-1 state for monopoles of different charges up to $5g_D$ [80]. The search result was interpreted by considering the Drell-Yan and photon-fusion mechanisms with mass limits ranging from 1500 to 3750 GeV. Insofar as proton-proton collisions are exploited, the MoEDAL Experiment sets stronger mass limits for charges of $3 - 5g_D$, while the ATLAS Experiment outperforms in the charge range of $1 - 2g_D$, as illustrated in Figure 1.5. The MoEDAL Collaboration started to show an interest in HECOs recently and reported the strongest Drell-Yan limits using Run-1 data for electric charges up to $|z| = 175$ [83]. As an example of extended search interpretations, with the Drell-Yan production mechanism, MoEDAL searched for dyons (particles with electric and magnetic charge) in proton-proton collisions, excluding magnetic charges up to $5g_D$ and

electric charges up to $|z| = 200$ for mass limits between 790 GeV and 3130 GeV [82].

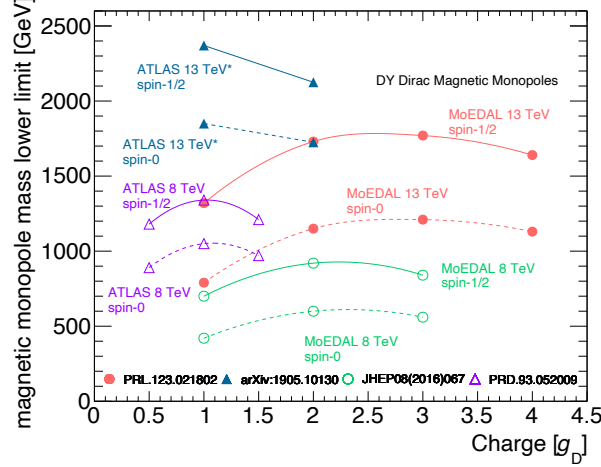


Figure 1.5: Comparison of monopole mass limits by ATLAS and MoEDAL searches in Run 1 and early Run 2 assuming the Drell-Yan (DY) production mechanism [74].

Searches in Heavy-Ion Collisions

In addition to collisions of elementary particles (electrons and positrons) or bound states of elementary particles (protons and antiprotons), HIP searches began to gain ground in heavy-ion collisions, which happened around the turn of the twenty-first century before the startup of the LHC. Previously, monopole searches were conducted at the Brookhaven Alternating Gradient Synchrotron using gold nuclei and at the CERN Super Proton Synchrotron using lead nuclei with lead targets, placing lower mass limits of 3.3 GeV and 8.1 GeV, respectively [85]. As far as HECOs are concerned, earlier results were presented by collaborations at the alternating-gradient synchrotron (AGS) of BNL [86–88], followed by a search using gold-gold collisions by the STAR Collaboration at the Relativistic Heavy Ion Collider (RHIC) [89]. Since the 1990s for a decade, these efforts provided model-independent cross-section limits for masses up to 100 GeV. The uniqueness of this collision system has been rediscovered in the recent past, leading to a MoEDAL search for monopoles with the LHC

heavy-ion collision data [81]. Utilizing the strong magnetic field created by the fast-moving heavy ions, it set an important precedent for search interpretations by assuming monopole production via the Schwinger mechanism [90–96]. As a result, the search managed to exclude charges $1 - 3g_D$ and masses up to 75 GeV.

1.3.2 Non-Collider Searches

Magnetic monopoles and HECOs are also targeted by non-collider searches. In particular, these searches look for GUT monopoles that might have been produced in the early universe [97], as well as nuclearites (nuggets of strange quark matter in collision with the Earth) that are either formed in the early universe or during collisions of neutron stars [35]. Predicted with a wide range of masses [6], the intermediate mass monopoles, i.e., with mass greater than 10^5 GeV, can be accelerated to relativistic velocities or higher by intergalactic magnetic fields, while heavier ones, i.e., with mass on the order of $\sim 10^{17}$ GeV, are primarily accelerated by gravitational interactions to nonrelativistic velocities. The lighter monopoles and nuclearites tend to have higher velocities; they can penetrate a large amount of material and thus are usually searched for in comic-ray detectors, while the heavy nonrelativistic monopoles are sought through catalysis of nucleon decay.

Numerous experiments have been devoted to searching for monopoles in cosmic rays [98–122]. Optimized for $1g_D$ monopoles, the MACRO Experiment is known for its sensitivity over the range $4 \times 10^{-5} < \beta < 1$ for the monopole’s normalized velocity β [102–104]. MACRO’s upper limits of the monopole flux on the order of $1.4 \times 10^{-16} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$ remain the best in the nonrelativistic regime for monopoles with mass above $\sim 10^{10}$ GeV [104]. Also worth noting are the SLIM Experiment, which searched for monopoles with mass in the range of 10^5 GeV to 10^{12} GeV in the nonrelativistic regime [107], as well as the NOvA Experiment, which probed mass down to 10^8 GeV for the same velocity range as MACRO’s [121, 122]. As for the other

side of the monopole velocity spectrum, the ANITA-II Experiment placed the strongest limits on the order of $1.0 \times 10^{-19} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$ on ultrarelativistic $1g_D$ monopoles [110]. Bridging the gap between the nonrelativistic and ultrarelativistic regimes, the IceCube Collaboration recently reported the best limit down to $2.0 \times 10^{-19} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$ on relativistic monopoles with $\beta \geq 0.8$ and $\beta \approx 0.995$ [120]. In addition to its sensitivity to monopoles in both the relativistic and ultrarelativistic regimes [115], the ANTARES Experiment is also interested in nuclearites [123, 124]. Competing against earlier searches [125, 126], the recent ANTARES search yielded the most stringent upper limits on nuclearites with mass above 10^{13} GeV and velocity $\beta \sim 10^{-3}$ [124]. In addition, the detection of monopole-induced nucleon decays by the Callan-Rubakov mechanism [127–129] is also used in searching for monopoles. Various searches reported β -dependent upper limits of flux in the range of $\sim 10^{-18} - \sim 10^{-14} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$ [130–135], while a search at Super-Kamiokande obtained $6.3 \times 10^{-24}(\beta/10^{-3})^2 \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$, the best upper limit for $\beta < 10^{-2}$ monopoles with mass 10^{16} GeV [136].

1.4 2015 – 2018 ATLAS Search

This work expounds a minimal-model based search for HIPs using 138 fb^{-1} of 13 TeV proton-proton collision data collected with the ATLAS Detector during the full Run 2 period of the LHC, with the result interpreted by the Drell-Yan and the photon-fusion production mechanisms. This chapter, Chapter 1, is a starting point for understanding the background to the full Run 2 ATLAS search. Next, Chapter 2 describes the benchmark models of this search, followed by their phenomenology. With a brief overview of the LHC and the ATLAS Detector in Chapter 3, the simulation of HIPs in the ATLAS setting is discussed in Chapter 4. Subsequently, Chapter 5 provides an account of the analysis strategy that leads up to the final result in Chapter 6. Chapter 7 envisions future prospects for searches in heavy-ion

collisions. Finally, Chapter 8 closes with conclusions and remarks.

Chapter 2

Collider Phenomenology of Highly Ionizing Particles

Monopole searches at colliders usually target stable Dirac magnetic monopoles and assume that magnetic charge is conserved [20, 23, 36, 37]. One conundrum with the Dirac monopoles is that their production cannot be handled with perturbation theory due to their strong magnetic couplings following the Dirac quantization condition, as per Equation 1.8. Consequently, there is no well-established theory of their production to date for direct cross-section calculation. However, in order to obtain a mass limit that can be compared to existing experimental results, it is necessary to assume at least one model at leading order with an associated production cross section. Hence, the common approach is to formulate a magnetic-dual theory of Quantum Electrodynamics (QED) and consider possible benchmark scenarios. The benchmark model for monopole production has been the Drell-Yan (DY) mechanism but the photon-fusion (PF) mechanism is now of interest, since at the current LHC energy scale its cross section in proton-proton collisions is larger [137–140]. In order to consider the photon-fusion mechanism for the analysis described in Chapter 5 and Chapter 6, I implemented this mechanism into a Monte Carlo event generator and then

validated the results. As an extension to the original Drell-Yan mechanism mediated by a virtual photon, I introduced and modelled the production of spin- $\frac{1}{2}$ HECOs mediated by a Z^0 boson. Furthermore, I used this generator to calculate the production cross sections and kinematic properties of the Drell-Yan and photon-fusion mechanisms at the Run 2 centre-of-mass collision energy. These studies have been published in [141], which however considers the production cross sections and kinematic properties of these production mechanisms at a higher centre-of-mass collision energy of 14 TeV.

These production mechanisms will be examined in this chapter. Section 2.1 identifies the interaction vertices for HIP production in DY and PF mechanisms. Section 2.2 covers the implementation of the interaction vertices. Section 2.3 presents production cross sections and representative kinematic distributions.

2.1 Partonic Mechanisms for HIP Pair Production

The approximate duality between electric and magnetic charges in electromagnetism is the key to the development of a magnetic-dual theory of QED, which is mediated by photon exchange. QED has an underlying invariance under $U(1)$ gauge transformations, where the field of each charged particle picks up a phase proportional to its electric charge. This invariance fixes the coupling of the field to the photon and demands electric charge conservation at each vertex. The electromagnetic coupling to the photon for any electrically charged particle scales as the charge $|z|$, where z is the electric charge in units of the elementary electric charge e . Hence, a minimal model of monopole interactions assumes an electromagnetic monopole-photon coupling that depends on the magnetic charge Ng_D of the monopole, where N is an integer. In the case of HECOs, it is completely natural to build additional interactions with a gauge invariance under $SU(2)$ transformations that govern the Standard-Model weak interactions, where the new gauge group gives rise to different matter-gauge couplings.

By constraining the possible interactions, the gauge symmetry allows me to write candidate Lagrangians from which I can compute the dynamics of the theory. I assume two spin states for the HIPs: spin-0 and spin- $\frac{1}{2}$. For a given spin, the candidate $U(1)$ Lagrangian is the same for monopoles and HECOs except for the coupling, which depends on Ng_D for monopoles and $|z|$ for HECOs. Hence, in the development below, I generalize monopoles and HECOs as HIPs with charge g . Working at leading order, I present two partonic HIP production mechanisms in proton-proton and heavy-ion collisions. The Drell-Yan mechanism is the process of a quark and an anti-quark annihilating to form a virtual photon (or a Z^0 boson in the case of HECOs), which then decays to a pair of HIPs. The photon-fusion mechanism is when two photons radiated from two colliding protons fuse to produce a pair of HIPs.

Invoking electric-magnetic duality, the charge and coupling of the magnetic monopole are often assumed to be velocity-dependent. This choice was based on the observed equivalence, in the small scattering angle limit, of the electron-monopole scattering differential cross section in [142] and the Rutherford scattering differential cross section, after substituting $\frac{g}{c}$ for $\frac{e}{v}$ [20] (further discussed in Section 4.1.1). The resulting magnetic charge $g_c^v \equiv g\beta$ not only appears in the description of monopole energy loss, but also when considering monopole pair production via the Drell-Yan and photon-fusion mechanisms [77–80, 137–140, 143]. In this thesis, I choose to follow the approach of a minimal model for Dirac monopoles where all HIP charges and couplings are assumed to be velocity-independent. Doing so leads to slightly different production cross sections and kinematic features. While the velocity dependence suppresses production of monopoles with lower velocities, a velocity-independent monopole coupling does not, thereby giving rise to softer kinematic distributions and higher production cross sections.

2.1.1 Spin-0 Interaction

The production of spin-0 HIPs in the magnetic-dual theory of scalar QED is described by the candidate Lagrangian

$$\mathcal{L} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} + (D^\mu\phi)^\dagger(D_\mu\phi) - m^2\phi^\dagger\phi, \quad (2.1)$$

where the covariant derivative

$$D_\mu = \partial_\mu + igA_\mu \quad (2.2)$$

couples the spin-0 HIP field, ϕ , to the photon field, A_μ , whose field strength is $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$. As a result, the covariant derivative component of Equation 2.1 gives rise to the HIP interactions

$$\mathcal{L}_{\text{int}} = \mathcal{L}_{3\text{pt}} + \mathcal{L}_{4\text{pt}}, \quad (2.3)$$

whose components

$$\mathcal{L}_{3\text{pt}} = -igA_\mu(\phi^\dagger\partial^\mu\phi - \phi\partial^\mu\phi^\dagger)$$

and

$$\mathcal{L}_{4\text{pt}} = g^2 A^\mu\phi^\dagger A_\mu\phi$$

dictate the vertices shown in Figure 2.1.



Figure 2.1: Two vertices given by scalar HIP-photon interactions.

While the Drell-Yan mechanism for scalar HIP pair production, shown in Figure 2.2,

realizes only the three-point vertex, the photon-fusion process, shown in Figure 2.3, has contributions from both the three-point and four-point vertices.

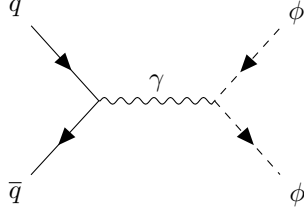


Figure 2.2: Feynman diagram of the photon-exchange Drell-Yan mechanism to produce a pair of scalar HIPs.

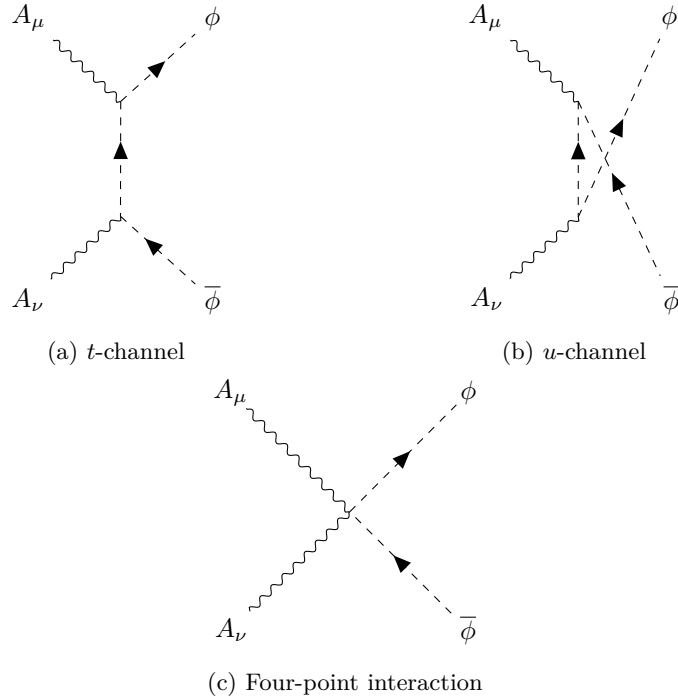


Figure 2.3: Feynman diagrams of photon fusion to produce a pair of scalar HIPs.

The interaction between the spin-0 HECOs and the Z^0 boson can be derived through the gauge invariance principle. Spin-0 HECOs couple to the Z^0 boson via quadra-linear interactions, therefore, they can only interact with a pair of Z^0 bosons assuming Standard-Model interactions, as shown in Figure 2.4. That is, the spin-0 HECOs cannot be produced

via the Z^0 -exchange-mediated Drell-Yan mechanism.

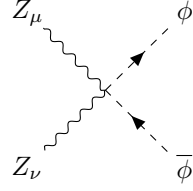


Figure 2.4: Feynman diagram of quadra-linear interactions of spin-0 HECOs with Z^0 bosons.

2.1.2 Spin- $\frac{1}{2}$ Interaction

Assuming HIPs are Dirac fermions, the basic Lagrangian for a spin- $\frac{1}{2}$ field ψ is

$$\mathcal{L} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} + \bar{\psi}i\gamma^\mu D_\mu\psi - m\bar{\psi}\psi,$$

with a different covariant derivative component of $\bar{\psi}i\gamma^\mu D_\mu\psi$ compared to that of the spin-0 scenario. This distinction gives rise to a different HIP-photon interaction term

$$\mathcal{L}_{\text{int}} = -g\bar{\psi}\gamma^\mu A_\mu\psi, \quad (2.4)$$

which corresponds to the three-point vertex of spin- $\frac{1}{2}$ HIPs in Figure 2.5. With this vertex,

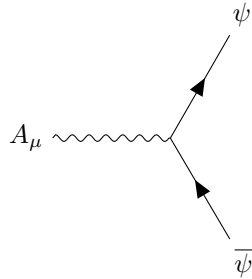


Figure 2.5: Three-point vertex for the interaction of spin- $\frac{1}{2}$ HIPs with a photon.

the photon-mediated Drell-Yan mechanism in Figure 2.6 was the first production process

considered for HIP production at the LHC [67, 70, 72, 77–80, 82, 83]. In addition, this three-point vertex contributes to the photon-fusion production process in Figure 2.7, where two such vertices are involved for a given production channel.

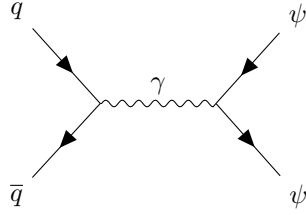


Figure 2.6: Feynman diagram of the photon-exchange Drell-Yan mechanism to produce a pair of fermionic HIPs.

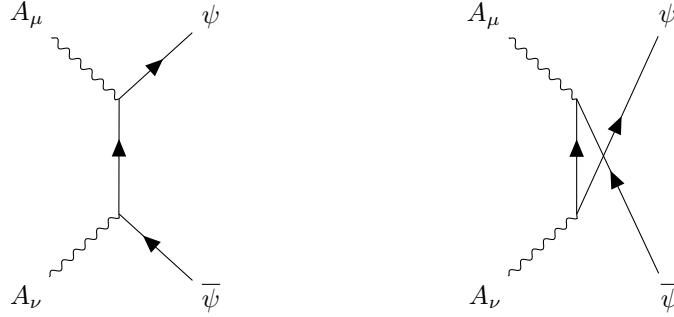


Figure 2.7: Feynman diagrams of photon fusion to produce a pair of fermionic HIPs.

While the magnetically charged monopoles are only allowed to interact via the electromagnetic neutral current, the electrically charged HECOs, like many other particles with exotic electric charges [144, 145], can also interact via the weak neutral current, as one assumes they transform under the $SU(2)$ symmetry of the weak interactions. To build such a model, I impose a local $SU(2)$ gauge invariance to generate a coupling to Z^0 bosons via a covariant derivative analogous to Equation 2.2. This interaction is given by

$$\mathcal{L}'_{\text{int}} = -\frac{e}{2 \sin \theta_W \cos \theta_W} \bar{\psi} \gamma^\mu (c_L P_L + c_R P_R) \psi Z_\mu, \quad (2.5)$$

where e is the electromagnetic coupling, θ_W is the weak mixing angle, and P_L and P_R are the

chiral projectors [146]. Unlike the HIP-photon interaction in Equation 2.4, this Lagrangian has a chirality dependence parametrized by two couplings that are generally not identical,

$$c_L = t_3 - |z| \sin^2 \theta_W \quad \text{and} \quad c_R = -|z| \sin^2 \theta_W, \quad (2.6)$$

where t_3 is the weak isospin of the HIPs. The assumption that HECOs transform as $SU(2)$ singlets implies that $t_3 = 0$ for HECOs. The interaction Lagrangian in Equation 2.5 is associated with the three-point vertex in Figure 2.8, which drives the Drell-Yan production of spin- $\frac{1}{2}$ HECOs via Z^0 -boson exchange, as in Figure 2.9.

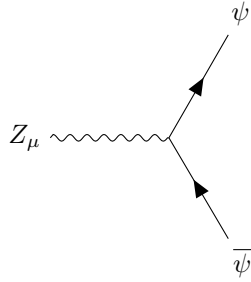


Figure 2.8: Three-point vertex for the interaction of spin- $\frac{1}{2}$ HECOs with a Z^0 boson.

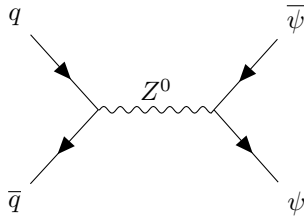


Figure 2.9: Feynman diagram of the Z^0 -exchange Drell-Yan mechanism to produce a pair of fermionic HECOs.

2.2 MadGraph Implementation

The HIP production model is built within MADGRAPH5_aMC@NLO [147], which is a generator that produces collision events. The model itself defines particles and their

interactions, which are passed to the MADGRAPH5_aMC@NLO generator for every user-specified process of interest. Next, MADGRAPH5_aMC@NLO searches for all interaction vertices in the specified model that accommodate the incoming and outgoing particles the user provides to generate a valid hard process through a set of Feynman diagrams. It then calculates the matrix element of the hard process from the Feynman diagrams via perturbation theory to a finite order. It calculates the production cross section as an integral of the matrix element over the initial momenta of the incoming particles and the phase space of the outgoing particles. In proton collisions, the probability of finding a parton, i.e., a constituent of the proton, at a given momentum follows the parton distribution function (PDF) [148]. The probability is distributed over a range of fractions x between 0 and 1, of the proton's momentum carried by the parton. The MADGRAPH5_aMC@NLO generator provides different PDF options via the Les Houches Accord PDF (LHAPDF) library [149].

2.2.1 Couplings

Every HIP interaction vertex described in Section 2.1 is characterized by a coupling constant that is proportional to the HIP charge. MADGRAPH5_aMC@NLO uses Heaviside-Lorentz units such that one electron charge is $\sqrt{4\pi}$ times the electron charge in Gaussian units. Thus, the fine structure constant is written as $\alpha = \frac{e^2}{4\pi}$, where the electron coupling e in Heaviside-Lorentz units is

$$e = 2\sqrt{\alpha\pi}. \quad (2.7)$$

As a result, the HECO couplings are integer multiples of the electron coupling, while the monopole couplings are integer multiples of g_D , obtained by rearranging Equation 1.10 and substituting Equation 2.7:

$$g_D = \frac{2\pi}{\sqrt{4\pi\alpha}} = \sqrt{\frac{\pi}{\alpha}}. \quad (2.8)$$

2.2.2 Spin-0 and Spin- $\frac{1}{2}$ Photon-HIP Vertices

The spin-0 HIP-photon interaction features the two vertices in Figure 2.1, where the three-point vertex is needed for both the Drell-Yan and the photon-fusion mechanisms, while the four-point vertex is only relevant for photon fusion. On the contrary, the spin- $\frac{1}{2}$ HIPs' interactions with a photon only feature the three-point vertex, which is responsible for both the Drell-Yan induced and the photon-fusion induced HIP pairs. For a given spin scenario, the Drell-Yan and photon-fusion mechanisms can be validated through the validation of the same vertices. For each vertex, it suffices to validate one interaction that involves it. Since the photon-fusion mechanism utilizes all three vertices, including the additional spin-0 four-point vertex, a successfully constructed photon-fusion mechanism guarantees the validity of its Drell-Yan counterpart. Hence, I validated the photon-HIP vertices in photon-fusion interactions.

Since the cross section depends explicitly on the centre-of-mass energy, it is ideal to work with collisions with a fixed centre-of-mass energy. However, the momenta of the photons emitted by two colliding protons follow a non-trivial distribution such that the effective photon collisions have no fixed centre-of-mass energy. For the ease of the theoretical calculations, I modeled the photon-fusion process in MADGRAPH5_aMC@NLO by colliding *bare photons*, i.e., photons that are not emitted by protons. In this case, the initial momenta of the photons are automatically set to a constant value given the fixed collision energy. I obtained cross sections from the photon collisions generated with MADGRAPH5_aMC@NLO. Subsequently, I calculated these cross sections analytically to validate the vertex implementation. Given all the vertices discussed in Section 2.1, the production cross sections at tree level are

$$\sigma(\gamma\gamma \rightarrow \phi\bar{\phi}) = \frac{g^4}{64\pi} \left[\frac{2\beta}{E^2} + \frac{2\beta m^4}{E^6} + \frac{E^2(1-\beta^2)^2 - 2m^2}{E^4} \ln \left(\frac{1+\beta}{1-\beta} \right) \right] \quad \text{for spin-0 HIPs,} \quad (2.9)$$

$$\sigma(\gamma\gamma \rightarrow \psi\bar{\psi}) = \frac{g^4}{32\pi E^2} \left[2\beta^3 - 4\beta + \frac{2E^4 + 2E^2m^2 - m^4}{E^4} \ln \left(\frac{1+\beta}{1-\beta} \right) \right] \quad \text{for spin-}\frac{1}{2} \text{ HIPs,} \quad (2.10)$$

where E is the energy of each photon, g stands for the coupling strength, β is the normalized velocity and m is the rest mass for each HIP.

When calculating photon-fusion cross sections with the MADGRAPH5_aMC@NLO model, I fixed the energy of each colliding photon to 6.5 TeV. Clearly, the 13 TeV centre-of-mass energy for the two-photon system is beyond the reach of the LHC. However, for validation purposes this energy can be arbitrarily chosen, as long as it is greater than the total mass of the final-state particles. I validated the MADGRAPH5_aMC@NLO model calculation against theory given by Equation 2.9 and Equation 2.10. For a fixed photon energy $E = 6.5$ TeV, the HIP normalized velocity β can be constrained by the conservation of energy as the HIP mass m and the square of the centre-of-mass energy of the incoming photons $s = (2E)^2$

$$\beta = \sqrt{1 - \frac{4m^2}{s}}. \quad (2.11)$$

The coupling strength g scales with the magnetic or electric charges the HIP carries, where the unit magnetic and electric charges are defined in Equation 2.8 and Equation 2.7 for magnetic monopoles and HECOs, respectively. Substituting E , m , β and g into Equation 2.9 and Equation 2.10, I obtained the theoretical cross-section values for $|g| = 1g_D$ monopoles in Table 2.1 and for $|z| = 60$ HECOs (for their approximate equivalence to $|g| = 1g_D$ monopoles dictated by Equation 1.10) in Table 2.2.

It should be noted that Equations 2.9 and 2.10 result from a calculation assuming an expansion of a perturbative process composed of a leading-order, or tree-level, contribution (reflected by the Feynman diagrams shown in Figures 2.3 and 2.7) and higher-order contributions of successively smaller corrections. However, due to the high HIP charge, the charge-squared-dependent HIP-photon coupling, i.e., 34 for $|g| = 1g_D$ monopoles, ultimately

Mass (GeV)	Spin-0 $ g = 1g_D$ monopole PF σ [pb]			Spin- $\frac{1}{2}$ $ g = 1g_D$ monopole PF σ [pb]		
	Model	Theory	Ratio (M/T)	Model	Theory	Ratio (M/T)
200	1.577×10^4	1.577×10^4	1.000	2.337×10^5	2.337×10^5	1.000
500	1.532×10^4	1.532×10^4	1.000	1.763×10^5	1.763×10^5	1.000
1000	1.417×10^4	1.417×10^4	1.000	1.343×10^5	1.343×10^5	1.000
1500	1.275×10^4	1.275×10^4	1.000	1.108×10^5	1.109×10^5	1.002
2000	1.126×10^4	1.125×10^4	1.001	9.485×10^4	9.488×10^4	1.000
2500	9832	9829	1.000	8.259×10^4	8.262×10^4	1.000
3000	8578	8579	1.000	7.230×10^4	7.233×10^4	1.000
4000	6858	6859	1.000	5.406×10^4	5.407×10^4	1.000
5000	6136	6144	0.999	3.578×10^4	3.578×10^4	0.999
6000	5023	5016	1.001	1.574×10^4	1.573×10^4	1.001

Table 2.1: Comparison of cross sections obtained from the model and calculated analytically for $|g| = 1g_D$ monopoles of various mass-spin combinations produced via photon fusion (PF) at a fixed $\gamma\gamma$ centre-of-mass energy of 13 TeV.

Mass (GeV)	Spin-0 $ z = 60$ HECO PF σ [pb]			Spin- $\frac{1}{2}$ $ z = 60$ HECO PF σ [pb]		
	Model	Theory	Ratio (M/T)	Model	Theory	Ratio (M/T)
200	1.061×10^4	1.061×10^4	1.000	1.572×10^5	1.572×10^5	1.000
500	1.030×10^4	1.031×10^4	0.999	1.186×10^5	1.186×10^5	1.000
1000	9.528×10^3	9.529×10^3	1.000	9.033×10^4	9.032×10^4	1.000
1500	8.575×10^3	8.573×10^3	1.000	7.453×10^4	7.457×10^4	0.999
2000	7.573×10^3	7.569×10^3	1.001	6.380×10^4	6.382×10^4	1.000
2500	6.613×10^3	6.611×10^3	1.000	5.556×10^4	5.557×10^4	1.000
3000	5.770×10^3	5.771×10^3	1.000	4.863×10^4	4.865×10^4	1.000
4000	4.613×10^3	4.614×10^3	1.000	3.636×10^4	3.637×10^4	1.000
5000	4.127×10^3	4.132×10^3	0.999	2.405×10^4	2.407×10^4	0.999
6000	3.379×10^3	3.374×10^3	1.001	1.059×10^4	1.058×10^4	1.001

Table 2.2: Comparison of cross sections obtained from the model and calculated analytically for $|z| = 60$ HECOs of various mass-spin combinations produced via photon fusion (PF) at a fixed $\gamma\gamma$ centre-of-mass energy of 13 TeV.

results in a divergence of the HIP production cross section in the perturbative expansion beyond leading order. Therefore, a non-perturbative approach is called for in order to study

the dynamics associated with the large coupling. Often used in the treatment of the low-energy regime of Quantum Chromodynamics (QCD), where the coupling between quarks and gluons increases beyond unity, one such approach constructs the theory on a finite cubic lattice of points representing the physical space-time coordinates so that the coupling is controlled by adjusting the lattice spacing. Since no one has attempted such an approach for the HIP cross-section calculation, I chose to neglect the higher-order corrections that do not converge and performed the tree-level calculation, which only serves as an approximation to a highly non-perturbative process.

2.2.3 Spin- $\frac{1}{2}$ Drell-Yan Z^0 -Exchange Vertex for HECOs

The implementation of the Z^0 boson decay to a pair of spin- $\frac{1}{2}$ HECOs is similar to that of the decay of a Z^0 boson to an electron-positron pair, apart from the necessary modifications to represent the zero weak isospin and $|z| > 1$ electric charges of the HECOs. I further examined this vertex with cross-section calculations in MADGRAPH5_aMC@NLO itself to look for the expected charge dependence. Again, the cross section was computed only with its tree-level contribution given by the Feynman diagram in Figure 2.9, so was the cross section of the photon-mediated Drell-Yan mechanism (based on the Feynman diagrams in Figures 2.2 and 2.6). The cross section σ goes as the absolute square of the matrix element $\mathcal{M}$, given by the interaction Lagrangian in Equation 2.5, which scales as the couplings $c_{L/R}$ in Equation 2.6:

$$\sigma \propto |\mathcal{M}|^2 \propto |c_L + c_R|^2 \propto |z|^2. \quad (2.12)$$

That is, the cross section of the Z^0 -exchange mode grows as $|z|^2$, just as in the photon exchange mode. Cross sections for both modes are listed for mass 200 GeV HECOs in Table 2.3 for five HECO charges centred at $|z| = 60$. The two exchange modes destructively interfere as a result of the relative sign between the two interactions parametrized in Equation 2.4

and Equation 2.5, hence, the total cross section is smaller than the photon-exchange cross section alone.

Charge $ z $	Drell-Yan σ for 200 GeV Spin- $\frac{1}{2}$ HECOs [pb]		
	γ exchange	Z^0 exchange	γ/Z^0 exchange
20	31.16	14.3	28.36
40	124.6	57.19	113.5
60	280.4	128.7	255.3
80	498.5	228.8	453.8
100	778.9	357.4	709.1

Table 2.3: Cross sections σ for mass 200 GeV spin- $\frac{1}{2}$ HECOs with varying charges in different Drell-Yan production modes.

2.2.4 Parton Distribution Functions

In proton-proton collisions, the incoming protons, denoted as A and B, are composed of partons, carrying initial momentum fractions x_i ($i = a, b$) with respect to the protons. The total cross section σ_{AB} of a collision between A and B can be expressed as the cross section $\hat{\sigma}_{ab}$ of parton collisions reweighted with the appropriate parton distribution functions (PDF) $f_{a/A}$ and $f_{b/B}$ for the initial state momentum fractions x_a and x_b :

$$\sigma_{AB} = \int dx_a dx_b f_{a/A}(x_a) f_{b/B}(x_b) \hat{\sigma}_{ab}. \quad (2.13)$$

If the partonic mechanisms consider different initial states, as for Drell-Yan and photon fusion, then the parton distribution functions $f_{a/A}$ and $f_{b/B}$ will be different. In MADGRAPH5_aMC@NLO, the NNPDF23LO_qed PDF [150] was used for Drell-Yan production while the LUXqed17 PDF [151, 152] was used for photon fusion, due to the intrinsic uncertainties of the PDFs. The quark and gluon PDFs are relatively well constrained in the NNPDF23LO_qed PDF, with uncertainties typically of order 5% [153], while the photon

PDF is weakly constrained due to the other processes present in the LHC collisions. With uncertainties up to 50% [154], the photon PDFs ultimately lead to large uncertainties in the predicted cross sections of the photon fusion process. This disadvantage is specifically targeted in the calculation of the LUXqed17 PDF, where electron-proton scattering data was collected to determine the photon PDF. The uncertainties in this measurement are significantly reduced to 1 – 2% [151] for a wide momentum range, as the momentum transfers in the scattering processes are carried by photons emitted from the protons.

2.3 Results

Production cross sections and kinematic distributions are compared for spin-0 and spin- $\frac{1}{2}$ HIPs for the LHC Run 2 collision energy of $\sqrt{s} = 13$ TeV. As per Tables 2.1 – 2.2, masses up to 4000 GeV are considered, limited by the collision energy.

The production cross sections of $|g| = 1g_D$ monopoles and $|z| = 60$ HECOs are shown in Figure 2.10, where different combinations of the partonic mechanisms and the spin states are presented. The photon-fusion mechanism is dominant at $\sqrt{s} = 13$ TeV. Also, the spin- $\frac{1}{2}$ HIPs have larger cross sections than the spin-0 HIPs, as explained below.

While the production cross sections are dependent on the production mechanism as well as the HIP spin, charge and mass, the kinematic distributions are dependent on all of these except for the HIP charge, as shown in Figure 2.11 and Figure 2.12 for a representative mass of 1500 GeV. The kinetic energy is softer for Drell-Yan production than for photon fusion and for spin- $\frac{1}{2}$ HIPs than for spin-0 HIPs. The pseudorapidity distributions are comparable in the two partonic mechanisms, but are more central for spin-0 HIPs than for spin- $\frac{1}{2}$ HIPs.

The difference in kinematic distributions between the Drell-Yan and the photon-fusion mechanisms arises from the difference in energies of the collision particles. While the energy of photons in the photon-fusion mechanism are limited only by the proton energy, the energy

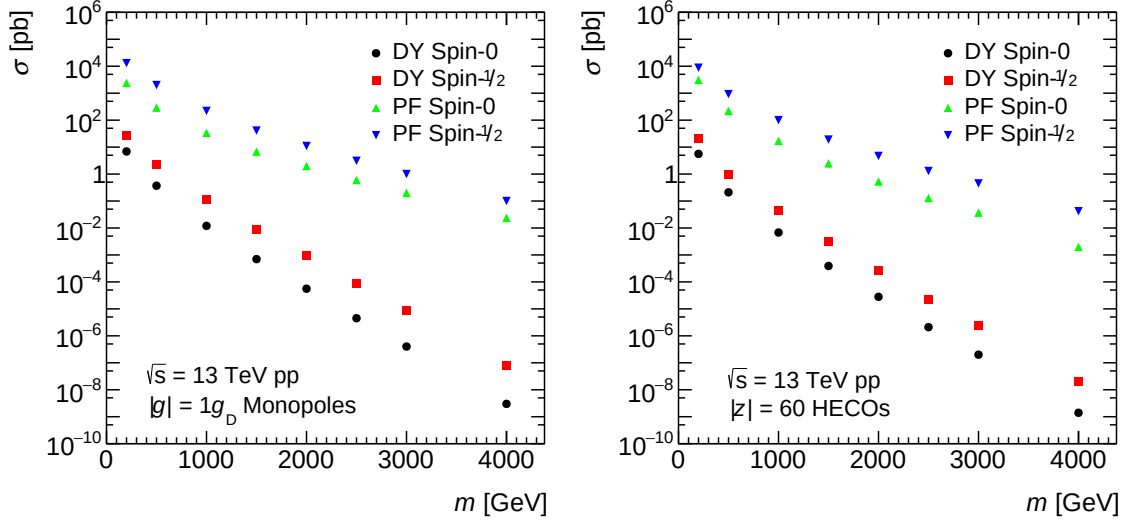


Figure 2.10: Cross sections for spin-0 and spin- $\frac{1}{2}$ HIPs via the Drell-Yan (DY) and photon-fusion (PF) mechanisms in proton-proton (pp) collisions.

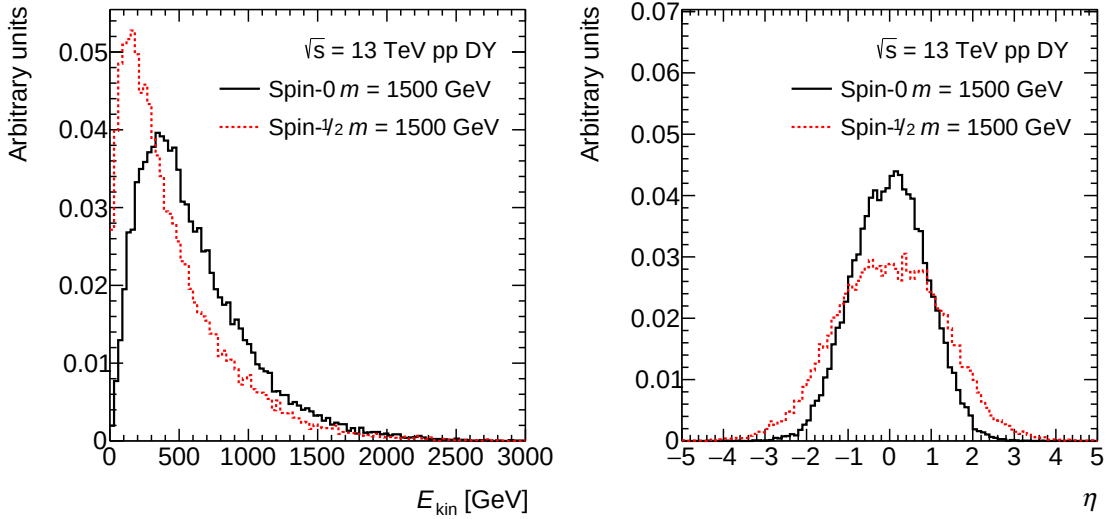


Figure 2.11: Without loss of generality, $|g| = 1g_D$ monopoles were chosen to show the kinetic energy (left) and pseudorapidity η (right) distributions of spin-0 (black solid line) and spin- $\frac{1}{2}$ (red dotted line) HIPs with mass 1500 GeV produced via the Drell-Yan (DY) mechanism in proton-proton (pp) collisions.

of one initial parton in the Drell-Yan mechanism has an additional constraint. As illustrated in Figure 2.2, Figure 2.6 and Figure 2.9, the Drell-Yan mechanism takes place when a quark (q) and an anti-quark ($\bar{q}$) collide. The quark can be an original constituent of the proton, but

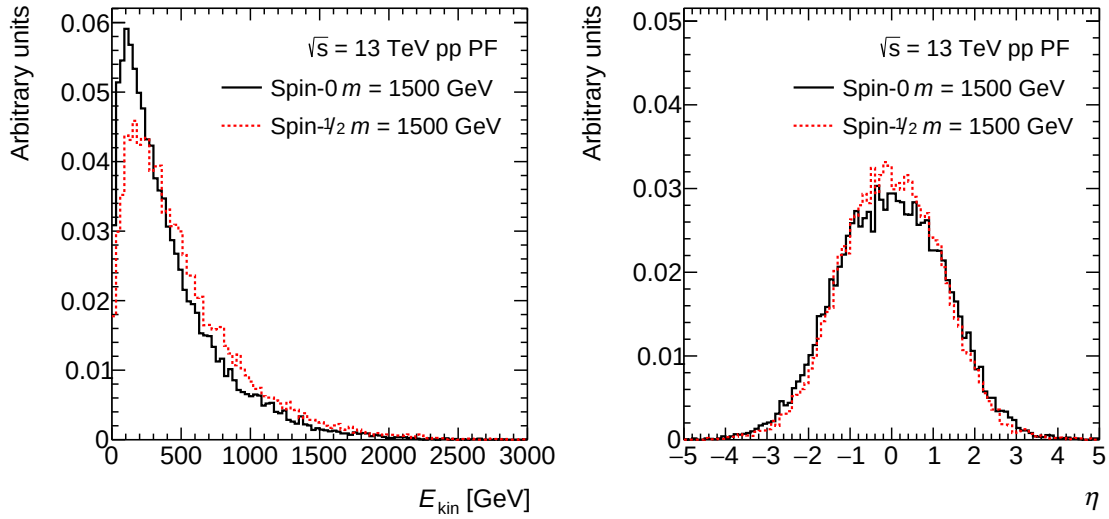


Figure 2.12: Without loss of generality, $|g| = 1g_D$ monopoles were chosen to show the kinetic energy (left) and pseudorapidity η (right) distributions of spin-0 (black solid line) and spin- $\frac{1}{2}$ (red dotted line) HIPs with mass 1500 GeV produced via the photon-fusion (PF) mechanism in proton-proton (pp) collisions.

the anti-quark must be converted from a gluon emitted by the proton in proton collisions. Therefore, the anti-quark typically has a lower energy than the quark, as well as the photon emitted directly from the proton. Consequently, the Drell-Yan mechanism has an overall lower energy than the photon-fusion mechanism, in terms of the colliding partons and the pair-produced HIPs. Therefore, HIPs of the same mass have higher kinetic energies when produced in the photon-fusion mechanism (see Figure 2.12) than the Drell-Yan mechanism (see Figure 2.11).

The difference in kinetic energy distributions between spin-0 and spin- $\frac{1}{2}$ HIPs needs to be discussed separately for the Drell-Yan and the photon-fusion mechanisms. In the Drell-Yan mechanism, the HIPs are produced from an intermediate boson (a photon or a Z^0 boson) of spin-1. By conservation of total angular momentum, i.e., the sum of spin and orbital angular momentum, spin-0 HIPs must be produced with non-zero orbital angular momentum, while spin- $\frac{1}{2}$ HIPs can be produced at rest without violating this conservation

law. Thus, spin-0 HIPs are produced with higher kinetic energies than spin- $\frac{1}{2}$ HIPs in the Drell-Yan mechanism. In the photon-fusion mechanism, the spin-0 HIP production can be realized via the four-point interaction (see Figure 2.3) while the spin- $\frac{1}{2}$ production only happens via the three-point interaction (see Figure 2.7). The four-point interaction in general produces HIPs of the same mass with lower kinetic energies than the three-point interaction, hence the lower kinetic energy in spin-0 production via the photon-fusion mechanism. The pseudorapidity distributions are correlated with the kinetic energy distributions, as high kinetic energy indicates higher Lorentz boost and thus lower pseudorapidity, as shown in Figure 2.11 and Figure 2.12.

2.4 Summary

In this chapter, I described the Drell-Yan and photon-fusion pair-production models of spin-0 and spin- $\frac{1}{2}$ magnetic monopoles and HECOs in proton-proton collisions. I implemented these models in MADGRAPH5_aMC@NLO to compute leading-order production cross sections and kinematic distributions in $\sqrt{s} = 13$ TeV pp collisions. These models will be used in the HIP search presented later in Chapter 5 and Chapter 6. Furthermore, I will extend the photon-fusion model for the production of HIPs in heavy-ion collisions in Chapter 7.

Chapter 3

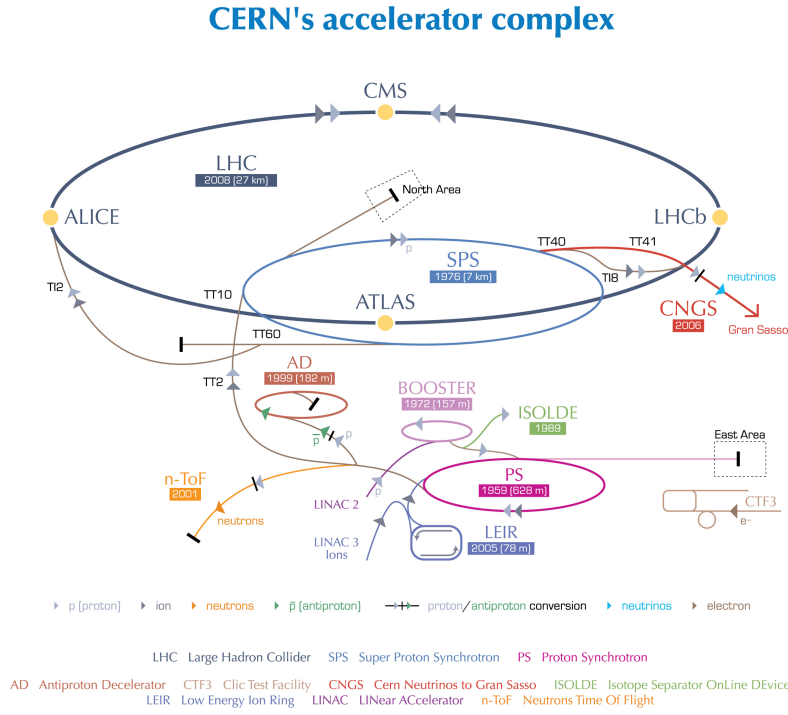
The ATLAS Experiment

As an understanding of the ATLAS Experiment is important for the Highly Ionizing Particle (HIP) analysis, key experimental components are briefly sketched in this chapter.

3.1 The Large Hadron Collider

The Large Hadron Collider (LHC) [155] delivers proton-proton collisions as well as heavy-ion collisions. The energies of these particles ramp up as they are passed through a series of accelerating facilities, as depicted in Figure 3.1. Protons, for example, are injected into Linear Accelerator 2 (Linac2) for a boost of energies to 50 MeV. Subsequently, they are sent through the Proton Synchrotron Booster, the Proton Synchrotron (PS) and the Super Proton Synchrotron (SPS), where they are accelerated to 1.4 GeV, 25 GeV and 450 GeV, respectively. Lastly, they are injected into the LHC ring, which accelerates the beams to their final energy (E) of 6.5 TeV in Run 2 (2015 – 2018). The acceleration is achieved by radio frequency cavities, whose oscillating electromagnetic fields push particles forward at a frequency of 400 MHz. Furthermore, particle beams are guided around the LHC ring as discrete bunches by a magnetic field from superconducting dipole magnets that keeps the

particles moving in circular trajectories. The preparation of heavy-ion collisions follows the same procedure to a large extent, except for an alternative linear accelerator connected to the ion source, and the Low Energy Ion Ring in lieu of the Proton Synchrotron Booster to provide shorter bunches. Circulating in the LHC ring, the beams with four-momenta $p = (E, \vec{p})$ and $p' = (E, -\vec{p})$ are focused by quadrupole magnets to attain the desired transverse cross-sectional sizes, as they are set to collide with a centre-of-mass energy of $\sqrt{s} = (p + p') = 2E$, at four interaction points, corresponding to four main detectors: ALICE, ATLAS, CMS, and LHCb. These experiments provide invaluable probes for new particles of high mass on the order of the centre-of-mass energies of the collisions.



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Figure 3.1: The CERN accelerator complex [156].

The second characteristic of the LHC is luminosity. The higher the luminosity, the better the statistical sensitivity of the measurements. The instantaneous luminosity, namely

the number of collision events per unit area per unit time, when summed over a period of time, gives the integrated luminosity, which is proportional to the total number of events delivered. This is used in conjunction with production cross sections to normalize data samples in physics analyses. For a fixed period of time, a gain of integrated luminosity is typically realized through the growth of instantaneous luminosity, by changing the number of particles per bunch, the number of bunches per beam, the transverse areas of the beams at the interaction points, and the frequency of collisions. During Run 2, proton-proton collisions are delivered with a frequency of 40 MHz, that is, bunches of protons collide, or cross, in the centre of a detector every 25 ns. Needless to say, the integrated luminosity will continue to accumulate as the LHC advances into further operational runs.

An inevitable consequence of a rise in instantaneous luminosity is the advent of pileup. *In-time pileup* refers to multiple inelastic interactions due to high-density particle beams in a single bunch crossing, whereas *out-of-time pileup* accounts for additional interactions from neighbouring collision bunches, or bunch crossings. Both types of pileup produce low-momentum particles that contaminate the signals of hard interactions and compromise the performance of event reconstruction. The number of inelastic interactions per bunch crossing can be extracted from the per-bunch instantaneous luminosity; it follows a Poisson distribution whose mean value μ is taken as a measure of pileup. A comparable measure is the number of inelastic interactions per bunch crossing averaged over a time block of events, denoted as $\langle\mu\rangle$. These two quantities are interchangeable if the number of interactions or the instantaneous luminosity does not change from bunch to bunch within the event block considered, but in reality one deviates from the other. Since $\langle\mu\rangle$ is measured over a longer time window than μ , it is more sensitive to out-of-time pileup that is related to the duration of the detector readout. In particular, if the detector reads out a longer time, say, 75 ns, than the bunch-crossing time of 25 ns, collision events that occurred in adjacent bunch crossings can be mistakenly attributed to the current bunch crossing. Also, the pileup can

be parametrized by tracing back these inelastic interactions to their locations of occurrence, known as primary vertices. Each interaction is associated with one primary vertex, thus, the number of primary vertices is another measure of the pileup, calculated for each bunch crossing as the number of reconstructed vertices, albeit vulnerable to inefficiency of the vertex reconstruction. All three quantities bring virtues to the estimation of pileup.

A summary of beam parameters and operating conditions is shown in Table 3.1 and Figure 3.2 for each year of data taking during Run 2. Commissioned in 2015, the Run 2 LHC launched collisions with 6.5 TeV per beam in a bunch crossing time of 25 ns. Over the next three years, the increase achieved in luminosities, through successively higher number of colliding bunch pairs and lower transverse size β^* of the quadrupole-focused beam at an interaction point, was astounding. Such changes brought about a peak instantaneous luminosity that exceeded the design value by a factor of two, and an integrated luminosity of 156 fb^{-1} delivered by 2018. Also notable is the pileup level, which went up significantly, especially from 2015 to 2016. Despite these challenges, the ATLAS Experiment operated successfully to utilize the LHC data for its physics program.

Parameter	2015	2016	2017	2018	Design
Energy [TeV]	6.5	6.5	6.5	6.5	7.0
Number of bunches	2232	2208	2554/1909	2544	2808
Bunch spacing [ns]	25	25	25/8b4e	25	25
Bunch population [10^{11} protons]	1.1	1.1	1.1/1.2	1.1	1.5
β^* [m]	0.8	0.4	0.3	0.3 - 0.25	0.55
Peak luminosity [$10^{34} \text{ cm}^{-2}\text{s}^{-1}$]	0.5	1.3	1.6	1.9	1.0
Peak pileup $\langle\mu\rangle$	16	41	45/60	55	-
Total integrated luminosity [fb^{-1}]	4.0	38.5	50.2	63.4	-

Table 3.1: LHC Run 2 parameters for proton-proton collisions at $\sqrt{s} = 13 \text{ TeV}$ during data-taking years 2015 – 2018, as well as their comparisons to the design values. The 2017 parameters are listed for two run configurations: the standard pattern of consecutive 25-ns bunches and a bunch structure of eight standard bunches followed by four empty bunches, denoted as “8b4e” [157].

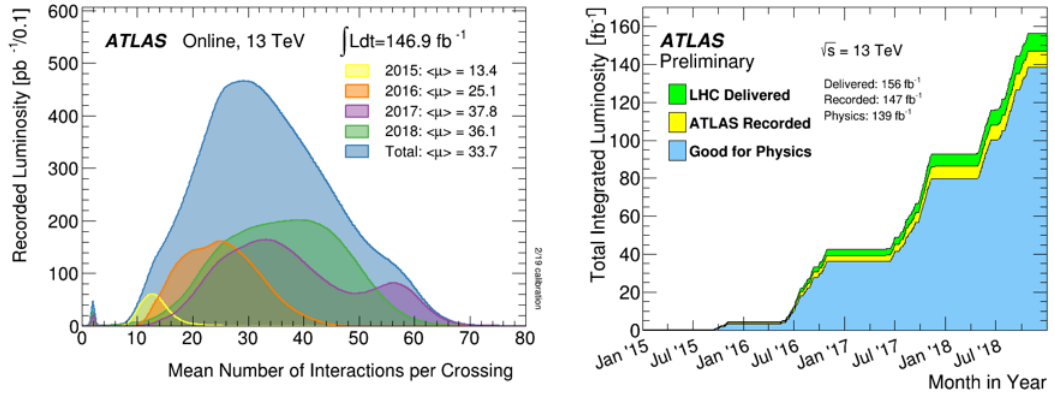


Figure 3.2: The mean number of interactions per bunch crossing (μ) as calculated from the ATLAS instantaneous luminosity with $\langle \mu \rangle$ denoting the average μ value for each data-taking year (left) and the ATLAS recorded integrated luminosity in comparison to the LHC delivered integrated luminosity (right) [158].

3.2 The ATLAS Detector

Built around the LHC beam pipe, ATLAS (A Toroidal LHC ApparatuS) is a cylindrically shaped general-purpose detector designed to measure trajectories and energies of particles produced in the hadronic collisions.

ATLAS uses a right-handed coordinate system with its origin at the nominal interaction point (IP) in the centre of the detector and the z -axis along the beam pipe. The x -axis points from the IP to the centre of the LHC ring, and the y -axis points upward. Cylindrical coordinates (r, ϕ) are used in the transverse plane, ϕ being the azimuthal angle around the z -axis. The pseudorapidity is defined in terms of the polar angle θ as $\eta = -\ln \tan(\theta/2)$ ¹.

Particles traverse concentric subsystems around the interaction point on their way out of the detector, whose main components are schematically drawn in Figure 3.3. Closest to the beam pipe, the Inner Detector (ID) operates within a 2-T axial solenoidal magnetic field to detect charged particles. Surrounding the ID are the calorimeters for measuring the energies of electrons, photons, and hadrons. The Muon Spectrometer, the outermost system,

¹This description of the ATLAS coordinate system definition is recommended by the ATLAS publications committee.

is designed to measure muon trajectories in a toroidal magnetic field. A trigger system is employed to identify these particles and record the associated collision events.

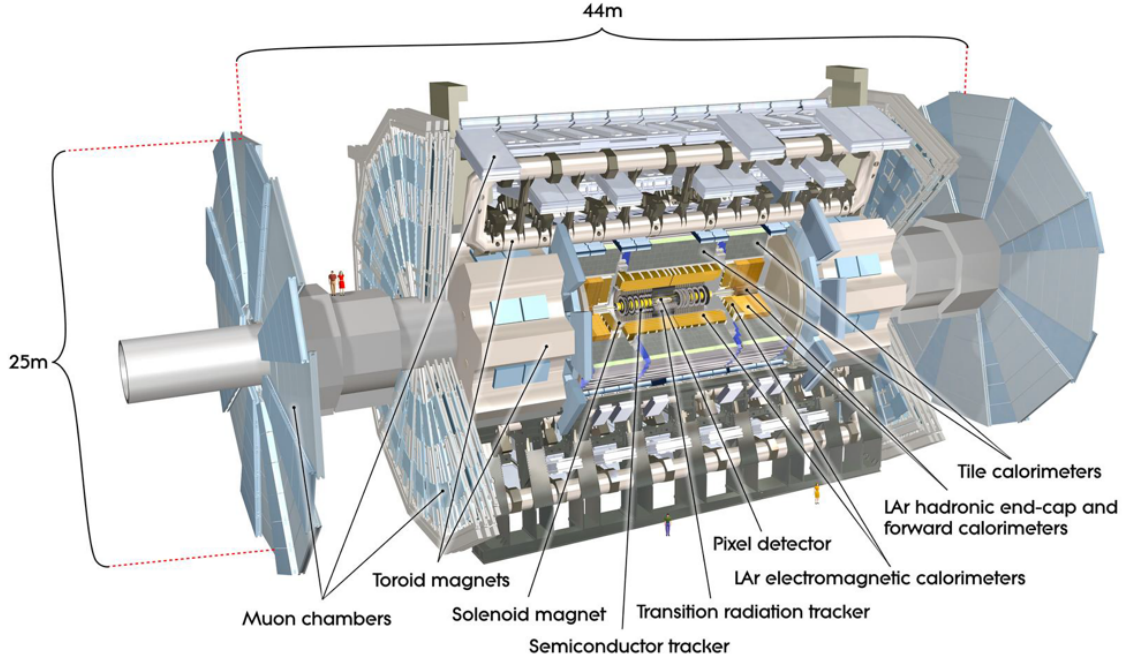


Figure 3.3: Cutaway view of the ATLAS Detector [159].

The collected data are processed into different streams and formats for worldwide distribution on the LHC Computing Grid (discussed in Section 3.3). In addition to the real collision data, there are simulated data that model the collision events and detector responses (discussed in Section 3.4).

3.2.1 Inner Detector

The Inner Detector (ID) consists of three subdetectors: the Pixel detector (including Insertable B-Layer), the SemiConductor Tracker (SCT), and the Transition Radiation Tracker (TRT), as shown in Figures 3.4 and 3.5. While the Pixel detector and SCT measure charged-particle trajectories (tracks) with high precision, the TRT features continuous tracking as

well as particle identification capabilities. The ID is separated into a barrel section oriented around the beam pipe that covers $|\eta| < 1.2$ and one end-cap perpendicular to the beam pipe on each side that extends the detector acceptance to $|\eta| < 2.5$. These subdetectors act on a binary read-out scheme, i.e., 1 is recorded if the signal in a sensor is above some threshold and zero recorded otherwise.

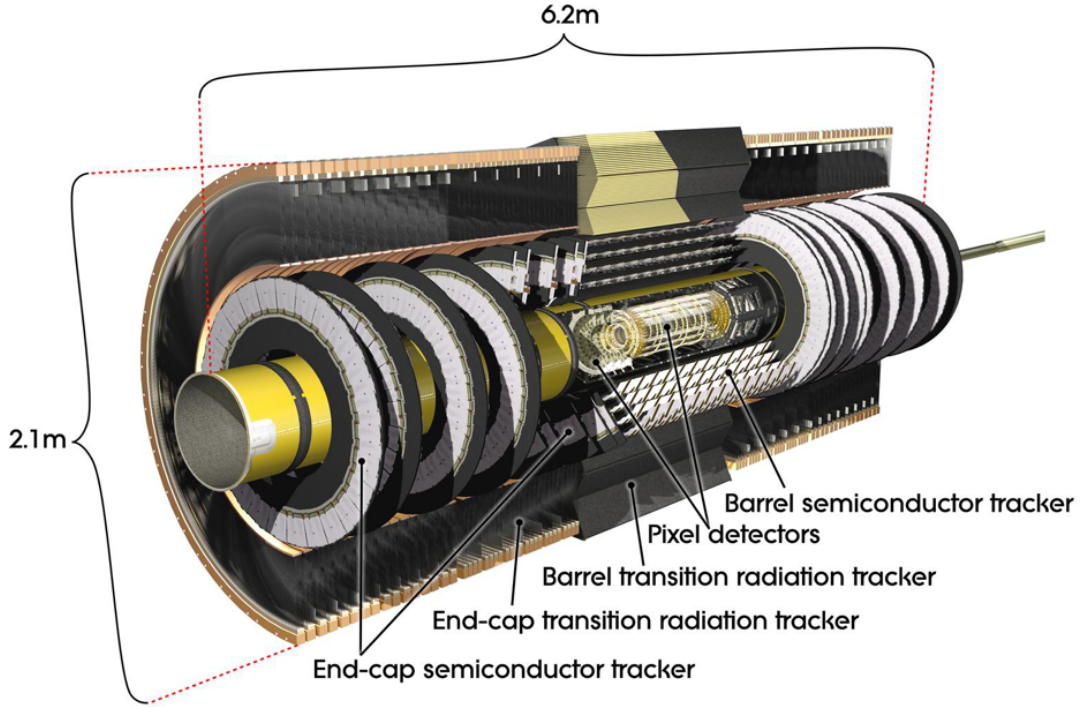


Figure 3.4: Cutaway view of the Inner Detector (ID) [159].

The Pixel Detector and IBL

As the innermost portion, the Pixel detector is composed of three concentric barrel layers and three disks in each end-cap. It capitalizes on silicon sensor technology, where charged particles traversing the detector lifts electrons from the valence band to the conduction band, thereby leaving holes in the valence band. The electron-hole pairs are separated by an electric field. As a result, the electrons from the electron-hole pairs enable a current

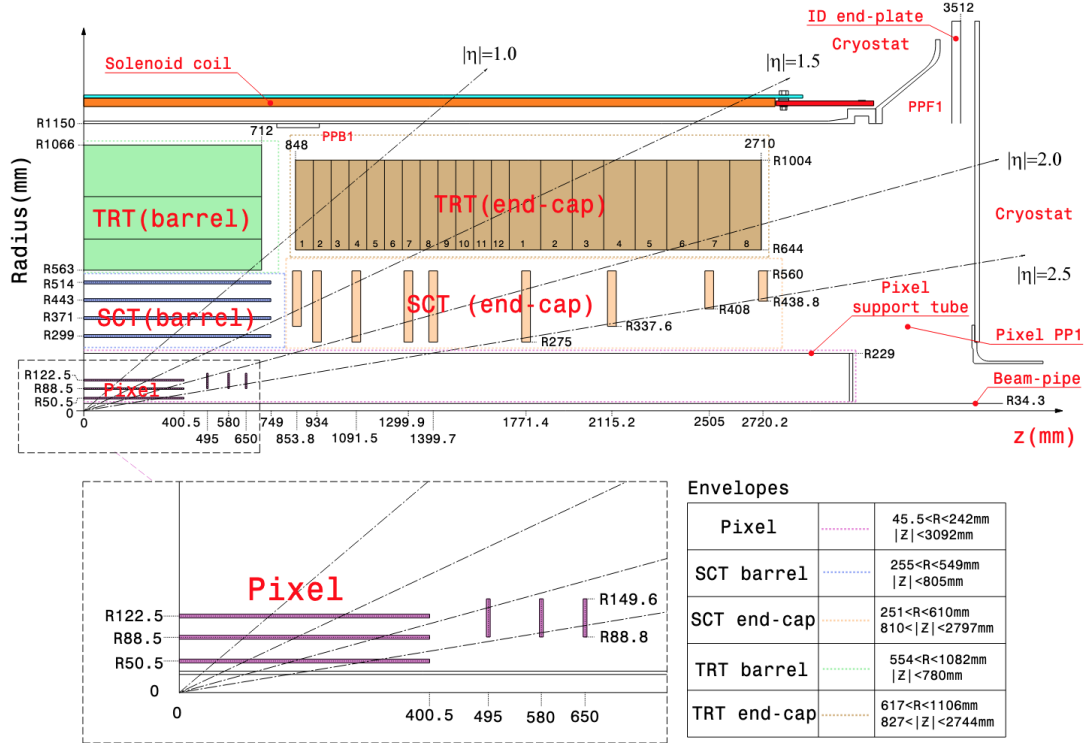


Figure 3.5: Transverse projection of the Inner Detector (ID) [159].

flow. Constructed with the finest granularity, it hosts a total of 80 million pixels with 1744 modules. Its high resolution, i.e., $(r - \phi) \times z = 10 \mu\text{m} \times 115 \mu\text{m}$ in the barrel and $(r - \phi) \times r = 10 \mu\text{m} \times 115 \mu\text{m}$ in the end-caps, allows for the measurement of impact parameters and the reconstruction of primary vertices and secondary decay vertices. Each pixel reads out a current pulse induced by ionization every 25 ns. The position (hit) of the pixel is recorded if the current pulse exceeds a threshold.

During the shutdown of the LHC in 2013 - 2014, an additional layer, i.e., the Insertable B-Layer (IBL) was installed in the barrel as the innermost layer. The IBL operates with readout chips and silicon sensor technologies. It maintains and improves the detector performance under high radiation and occupancy with a resolution of $(r - \phi) \times z = 8 \mu\text{m} \times 40 \mu\text{m}$.

SemiConductor Tracker

The SemiConductor Tracker is situated between the Pixel tracker and Transition Radiation Tracker. The detection of charged particles is realized with a technology as that of the Pixel detector and allows for precise tracking. It comprises 4088 double-sided modules, distributed over a four-layer barrel and one end-cap per side with nine disks each. The modules are segmented into six million silicon strips. It achieves a resolution of $(r - \phi) \times z = 17 \mu\text{m} \times 580 \mu\text{m}$ in the barrel and a resolution of $(r - \phi) \times r = 17 \mu\text{m} \times 580 \mu\text{m}$ in the end-caps.

Transition Radiation Tracker

Extending to 95 cm in radius, the Transition Radiation Tracker [160] constitutes the outermost layer of the ID. It was designed with two main purposes: particle tracking through the collection of ionization electrons in the active gas and electron-pion discrimination through the detection of transition radiation (TR) photons. It is instrumented with drift tubes (straws) made of aluminized kapton, coated by carbon polyimide and laminated with carbon fiber. The straws are filled with gas and are inserted with anode wires in their cavities. They originally operated with a xenon-based gas mixture, but some now utilize an argon-based mixture due to an increasing number of leaks developed over time resulting in a cost issue, as xenon gas is expensive. When charged particles pass through the gas in the straws, they produce ionization electrons, which drift in an electric potential of about -1500 V towards the central wires that are grounded. These electrons initiate ionization cascades, which give rise to current pulses read out at the ends of the wires. A total of 245,760 4-mm diameter straws are assembled into 73 and 160 layers for the barrel and end-caps, respectively. This configuration ensures that a particle with $|\eta| \leq 2.0$ will cross more than 30 straws and produce more than 30 hits with a resolution of 130 μm per hit, greatly enhancing the tracking

capability of the ID. Furthermore, the straw layers are interleaved with radiators for the production of transition radiation. The radiators are made of polypropylene fibers in the barrel and foils in the end-caps. As a charged particle crosses the boundary between the radiators and a straw, it experiences a change of index of refraction and emits photons in its direction of motion at an intensity proportional to its Lorentz factor γ . Given that this effect occurs at lower energies for lighter particles, electrons produce more TR photons than pions of the same energy. The TR photons are subsequently detected via photoelectric absorption, whose cross section scales as Z^5 with Z being the atomic number of the absorber. Thus, the TRT operates with high- Z straw gas and low- Z radiators.

The signal is read out by front-end electronics in two steps. First, a front-end chip called the Amplifier-Shaper-Discriminator Baseline Restorer (ASDBLR) amplifies, shapes and discriminates the readout signal against two thresholds: low threshold (LT) and high threshold (HT). The LT discriminates against ionization and δ -rays (high-energy electrons liberated from inner orbitals of atoms) that typically deposit 2 keV, thus is used to measure a particle's position. Within the range of 100 to 300 eV, the LT was set to vary so that the noise occupancy (rate) is uniform across the TRT [161]. On the other hand, the HT is used in particle identification based on the energy deposition by TR photons in the range of 5 to 15 keV. Since the photons are more readily absorbed in the xenon mixture than in the argon mixtures, the argon HT is lower, with its exact value tuned so that minimum ionizing muons have the same probability of being recorded as HT hits in both gas mixtures. While the HT is set to 5 – 6 keV in the xenon straws, it is set to 2 – 3 keV in argon straws. The variation in the HT setting ensures that the rejection of pions is uniform across the detector [162]. Subsequently, another front-end chip called the Drift-Time Measuring Readout Chip (DTMROC) measures the timing of the signals from the ASDBLR. The signal is processed in multiple equal-time bins during a 75-ns readout window (see Figure 3.6). For every 3.125 ns bin, the signal is digitized into a bit of one if the pulse exceeds the LT and zero

otherwise. The duration of the pulse is characterized by the number of time bins with non-zero bits, known as the time-over-threshold. Likewise, the energy deposition is compared against the HT every 25 ns, generating three binary HT bits in total during the readout window.

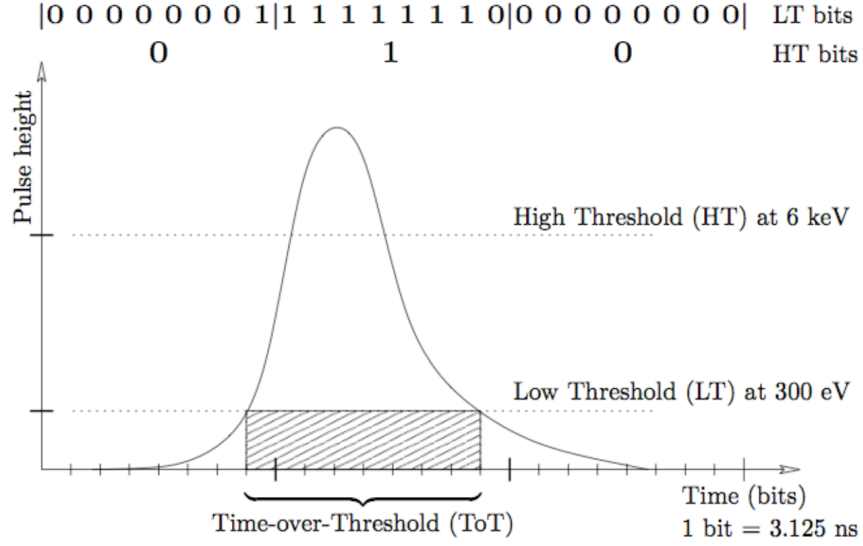


Figure 3.6: A typical pulse during the 75 ns readout window [163].

3.2.2 Calorimeters

The calorimeters are built from alternating layers of metal absorbers and active materials, such that when an incident particle produces a shower of energy, only a fraction of that energy is sampled by the active medium, and the rest is lost to the absorbers. This process of sampling, where the energy of the full shower is not directly measured but inferred from the energy of the observed shower, underlies the detection of electrons, photons, and hadrons up to $|\eta| = 4.9$. A schematic of the calorimeters is shown in Figure 3.7.

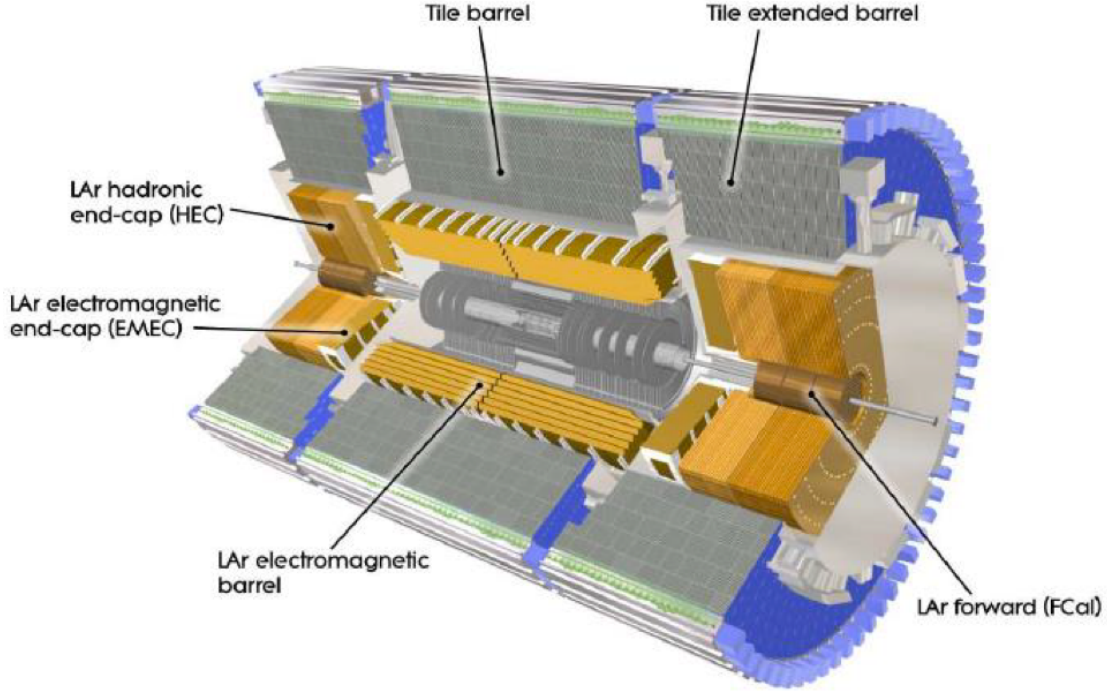


Figure 3.7: Schematic diagram of the calorimeters [159].

Electromagnetic Calorimeter

The Liquid Argon (LAr) Electromagnetic (EM) Calorimeter measures the electromagnetic showers of photons and electrons with a LAr-based active medium. A complete and uniform coverage in ϕ is supported by its accordion geometry, shown in Figure 3.8, while an extensive coverage in η results from the combination of a barrel section in the range $|\eta| < 1.45$ and two end-caps in $1.45 < |\eta| < 3.2$. A three-layer electrode is interleaved between every two consecutive lead absorbers. The two outer layers are connected to a high-voltage supply, which generates an electric field across the LAr gap between each absorber-electrode pair. The electric field drifts the ionization electrons towards the electrode, whose central layer reads out the signals from the gaps. The LAr EM Calorimeter is instrumented with four azimuthal readout layers, each of which contains readout cells formed by etching electrodes in $\eta \times \phi$ (see Figure 3.8). The readout granularity of each layer is set by the cell size $d\eta \times d\phi$,

which ultimately follows from the electrode etching pattern. The first compartment (EM1), known as strips, integrates cells of $d\eta \times d\phi = 0.003 \times 0.1$, achieving the finest granularity in η among all layers. The second sampling (EM2) is deeper and thus contains the majority of the energy shower. It is highly segmented in both directions with a readout granularity of $d\eta \times d\phi = 0.025 \times 0.025$. If the shower is not fully contained in the second sampling, the third sampling (EM3) adds sampling depth with a coarser division of $d\eta \times d\phi = 0.05 \times 0.025$. To account for energy loss before the EM Calorimeter, a presampler is placed in front of the EM1 layer in $|\eta| < 1.8$ to sample the early stage of the shower cascade with $d\eta \times d\phi = 0.025 \times 0.1$ per cell. Between the barrel and each end-cap, within $1.375 < |\eta| < 1.52$, is a transition region, which can be avoided as appropriate in analyses as it contains a large quantity of inactive material. This material, consisting of cables and pipes, provides services to the Inner Detector yet cannot accurately measure the energies of incident particles.

The ionization electrons liberated from the argon atoms drift to the copper electrodes within approximately 450 ns. The electrodes read out a triangularly shaped signal that peaks shortly upon its arrival and soon vanishes linearly. A shaping function is applied to the raw signal, producing a bipolar pulse with a much shorter readout time of 25 ns. As per Figure 3.9, signals from the current collision bunch give rise to the peak of the bipolar pulse, whereas those before or after this bunch crossing reduce the pulse. This shaping process not only enables triggering but separates in-time pileup from out-of-time pileup; it leaves the in-time pileup signals positive and smears out the out-of-time pileup signals, which can possibly be made negative.

Hadronic Calorimeters

The hadronic calorimeters, as their names suggest, predominantly measure hadronic showers. One component is the Tile hadronic calorimeter (Tile), which uses scintillator tiles as the active material and steel as the absorber in a barrel and two extended barrels.

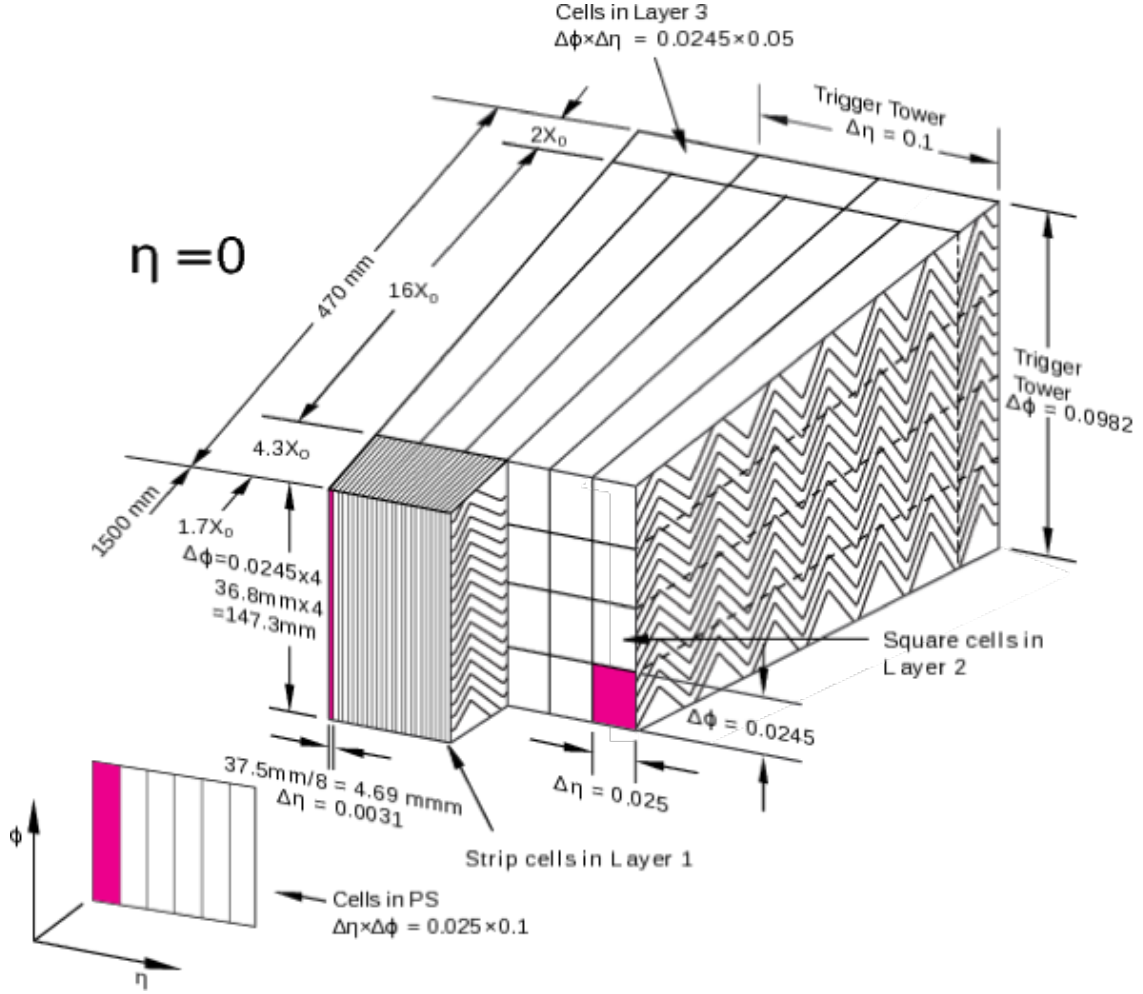


Figure 3.8: A section of the electromagnetic barrel calorimeter [159].

Compared to the Electromagnetic (EM) calorimeters, the Tile calorimeter has a narrower readout window of 150 ns and a coarser granularity with cells of 0.1×0.1 or 0.2×0.1 in $\eta \times \phi$. Located behind the EM end-cap calorimeter, each Hadronic End-cap Calorimeter (HEC) consists of two wheels instrumented with copper-plate absorbers in LAr gaps and reads out signals every 450 ns with varying $\eta \times \phi$ granularities of 0.1×0.1 and 0.2×0.2 .

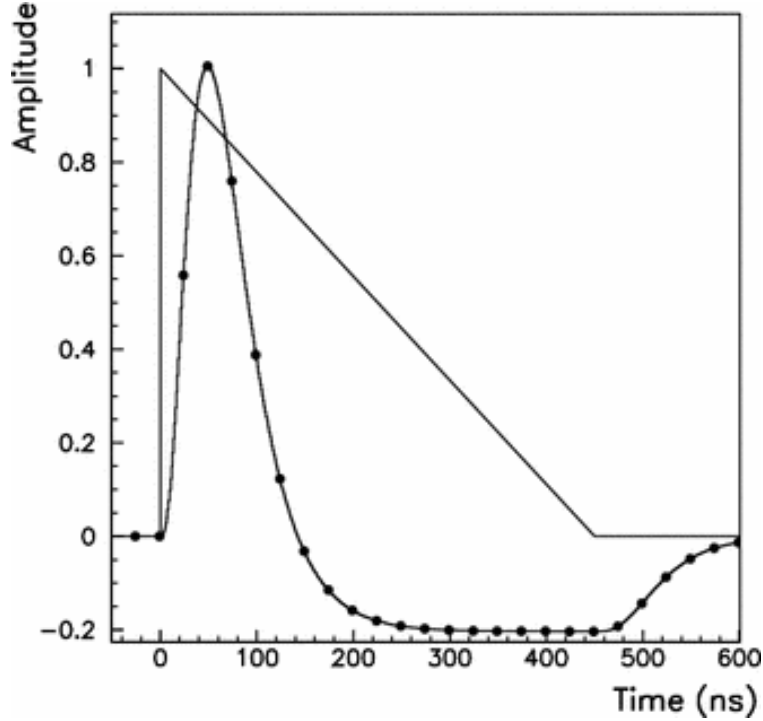


Figure 3.9: LAr EM Calorimeter current pulse in the detector and from the front-end electronics after shaping. The dots are sampling points every 25 ns [159].

Forward Calorimeter

Integrated into the end-cap calorimeters and capable of measuring both electromagnetic (EM) and hadronic showers, the Forward Calorimeter (FCal) extends the ATLAS coverage to close to one degree from the beams at the expense of exposure to intense radiation. This challenge motivates the design of a high-density metal matrix surrounding concentric electrodes, each of which holds a solid metal rod within a copper tube and LAr between the rod and the tube. The FCAL employs a relatively fast readout of 150 ns to reduce the effect of pileup on its signal pulse. Among three modules in each end-cap, only the first module is assembled with copper in the rods and in the metal matrix for optimized EM measurements. The other two modules of the FCal in each end-cap are optimized to measure hadronic showers with a different metal – tungsten, in the rods and for the metal

matrix, despite being similar in structure to the first module.

Calibration

The signal measurements need to be calibrated. The time calibration synchronizes the cell readout time to the LHC bunch crossing time, thereby assigning energy deposits to the correct bunch crossings. In the LAr Electromagnetic (EM) Calorimeter, for instance, the conversion of ionization current to energy deposited in the LAr gaps is calibrated to the *electromagnetic energy scale* that fully reconstructs the energy of an EM signal. Additionally, the amount of energy lost in the lead absorbers is accounted for to reflect the energy loss in the passive materials. Also, an integral part of the energy calibration is the handling of pedestals and noise. Monitored in empty bunch crossings where no collisions are expected, the pedestal is the baseline added to the signal such that the pulse is above the electronic noise and that the pulse amplitude yields a non-negative energy. The pedestal value for each cell is computed as an average over multiple events, whereas the noise is computed as the root-mean-square of the pedestal value. They are applied to cells, which then can be positive or negative in energy due to fluctuations around these values.

3.2.3 Muon Spectrometer

The Muon Spectrometer (MS) measures the positions and momenta of muons from their deflections in a toroidal magnetic field via two high-precision tracking chambers: Monitored Drift Tubes (MDTs) for the majority of the MS and Cathode Strip Chambers (CSCs) in the forward region. The MS provides the robustness needed to measure these particles with an axial (z) and radial (r) resolution of about 40 μm . The maximum charge collection time of 700 ns for the MDTs and 40 ns for the CSCs makes these two chambers ill-suited for triggering purposes, which however are fulfilled with Resistive Plate Chambers (RPCs) in

the barrel region and Thin Gap Chambers (TGCs) in the end-caps. These trigger chambers have a fast readout time no longer than 25 ns at the cost of worsened position resolution on the order of 1 mm.

3.2.4 Trigger System

The Trigger and Data Acquisition system (TDAQ) [164] transmits collision data from the front-end electronics to a two-tier trigger system via Readout Drivers (ROD) (see Figure 3.10). The trigger selection starts with hardware-based First-Level (L1) triggers that reconstruct energy deposits in the calorimeters and hits in the Muon Spectrometer (MS). The L1 trigger objects are stored in buffers and used as seeds to search for objects in smaller detector regions, known as regions of interest (RoIs). Software-based High-Level Triggers (HLT) initiate full-detector readout and decide whether an event should be saved or discarded based on detected patterns. The L1 accepts at most one per 400 events, reducing the rate of 40 MHz dictated by the LHC bunch-crossing rate to a maximum of 100 kHz. Subsequently, the HLT further refines the L1 decisions for permanent storage at an event rate of 1.5 kHz and a data rate of 1.5 GB/s. The L1 and HLT rates can be adjusted by prescale settings to select a fraction of the events that pass the corresponding triggers.

L1 Trigger

The goal of L1 is to identify isolated high-energy objects in the $\eta - \phi$ transverse planes of the full calorimeters and Muon Spectrometer.

The calorimeters trigger on electrons, photons, taus, and jets via a preprocessor, as well as two parallel processors, i.e., Cluster and Jet/Energy sum. The preprocessor digitizes and calibrates the analogue signals of the calorimeters. The Cluster processor identifies electron, photon or tau objects in windows of size $d\eta \times d\phi = 0.1 \times 0.1$, as discussed below. Cell

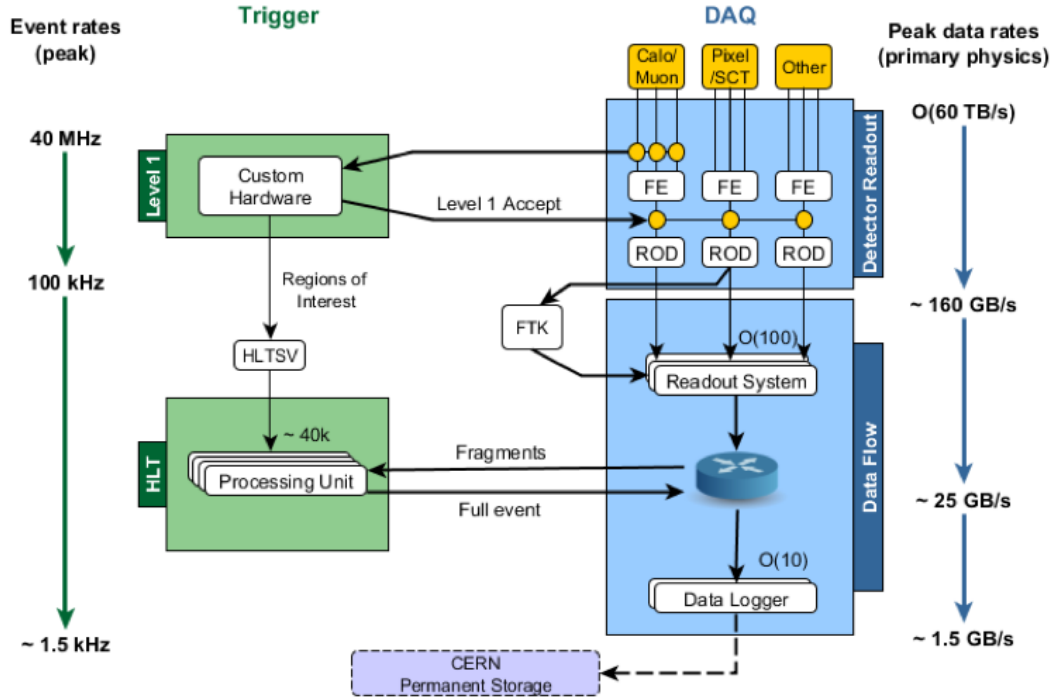


Figure 3.10: Schematic diagram of the ATLAS TDAQ system with expected peak rates and bandwidths [165].

signals are aggregated into towers of this granularity, which in turn are processed by the L1 trigger to form 4×4 tower clusters in the electromagnetic (EM) and hadronic calorimeters, as shown in Figure 3.11. Each trigger tower is associated with a transverse energy value E_T , defined as $E_T = E \sin \theta$, where E is the deposited energy and θ is the angle between the beam axis and a line from the interaction point to the centre of the cell. Hence, different combinations of trigger towers within each tower cluster give rises to different E_T sums. One important substructure embedded in a tower cluster is a(n) EM (hadronic) inner core, built from four central towers in a 2×2 window in the EM (hadronic) calorimeter. From the fact that each tower can be paired with another either horizontally or vertically, it follows that a total of 16 E_T sums can be attained in the full calorimeters by adding EM tower

pairs in the EM inner core to hadronic tower pairs in the hadronic inner core. An RoI is defined as a 2×2 tower cluster summed in depth over the two calorimeters if it contains a local E_T maximum and at least one E_T sum that exceeds a given threshold. Typically, other requirements are imposed on the RoI for it to be triggered, such as isolation thresholds on isolation rings formed by the 12 towers around the EM (hadronic) inner core, as well as veto thresholds on the hadronic inner cores for triggering on electrons and photons. These requirements are waived for EM clusters with transverse energies above 50 GeV. Lastly, the Jet/Energy processor identifies jet objects with sliding windows of $d\eta \times d\phi$ from 0.2×0.2 to 0.4×0.4 . It produces RoIs as 2×2 to 8×8 tower clusters if their E_T sums over the electromagnetic and hadronic calorimeters exceed a threshold.

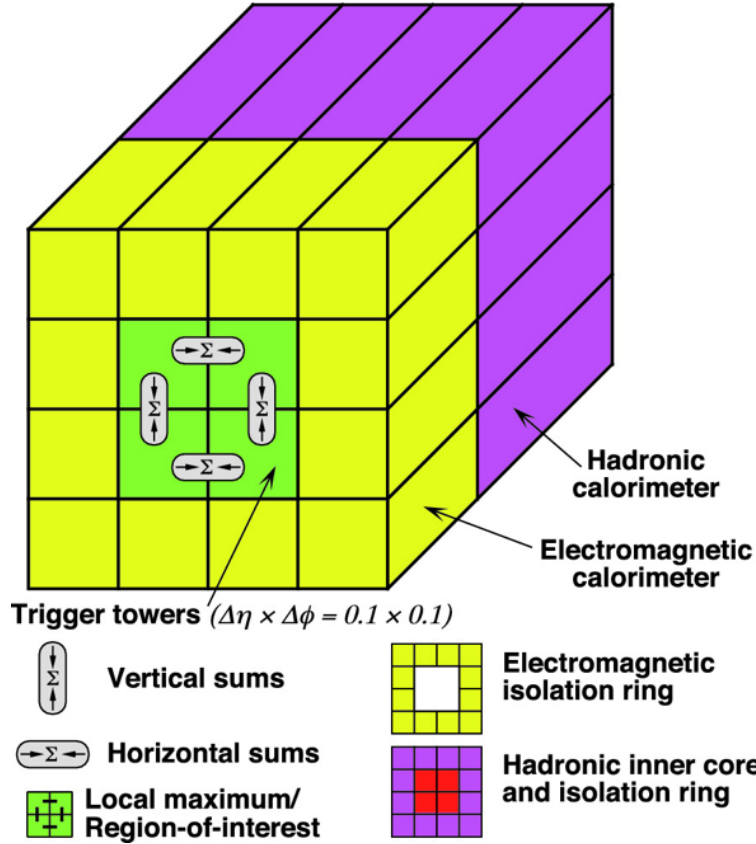


Figure 3.11: Trigger towers processed by L1 calorimeter triggers [159].

The L1 muon triggers recognize patterns of hits from the RPCs in the barrel and TGCs in the end-caps to select muons with a range of transverse momenta p_T . As a high- p_T muon tends to produce hits concentrated in a narrow road, muons in different p_T regimes are selected by different thresholds applied to their manifested roads.

HLT Trigger

The HLT makes more informed decisions as it performs reconstruction and identification not only on L1 RoIs but also on information not available at the L1, namely full-granularity calorimeter readout, precision measurement by the Muon Spectrometer and tracking information from the Inner Detector (ID). Furthermore, driven by latency requirements, the HLT follows a two-stage approach, where a fast first pass is executed to filter out unwanted events and leave the remaining for precision reconstruction.

Electron and photon candidates are reconstructed as clusters of cells within the L1 EM calorimeter RoIs. The second layer of the EM calorimeter is where electrons and photons deposit most of their energy and thus is used to locate cells with the largest transverse energies in the first pass. Fast tracking is done on the ID for the electron candidates. The separation in η between a reconstructed track and a cluster cannot be greater than some cut value for them to be correlated. In the detailed precision step, the cells are clustered by algorithms if their signals are compatible with the expected energy flow of a shower cascade. The sliding-window algorithm [166] searches for a region in the EM calorimeter whose total energy is a maximum, whereas the topological clustering algorithm gathers signal cells due to electromagnetic and hadronic showers by rejecting noise cells. The precision tracking associates electron tracks to clusters using the second layer of the EM calorimeter alongside a tighter cut on their separation.

Despite the improvement with respect to L1, the HLT neither processes full event data nor applies full calibration over the entire detector. Performed in a different stage, these are

separate procedures that will be addressed in Section 3.3.2, and thus are only mentioned here in passing.

3.3 Collision Data

The ATLAS Detector and its subsystems are on standby as soon as the LHC declares stable beams. They work alongside each other to record data that are subsequently processed into different streams for system calibrations and physics studies.

3.3.1 Data Taking

Each year of data collection progresses in successive runs, which are assigned unique numbers and grouped into alphabetically labelled *periods* if stable conditions are shared among them. A data-taking run corresponds to a single fill of the LHC whose beam luminosity is maximal at the start and gradually degrades over time, typically from a few hours to a day. Each run is divided into *luminosity blocks* of approximately one-minute time intervals where the instantaneous luminosity is more or less constant. Each luminosity block is also characterized by constant detector and trigger configuration as well as data quality conditions. The luminosity blocks that pass a set of data quality checks enter *Good Runs Lists*, which are made available for analysis purposes.

3.3.2 Data Processing

The RAW data output in the byte-stream format from the HLT consists of main streams such as the primary physics stream of all physics events, the express stream for a subset of events, or the calibration stream of partial detector events, e.g., in the ID. They are transferred to the CERN Tier-0 (T0) data centre – the first tier of the Worldwide LHC

Computing Grid (WLCG), which provisions reconstruction of the calibration and express streams, first-pass processing of the primary event stream, and final distribution of the derived datasets to the Tier-1 (T1), as described in Section 3.5. For example, in the calibration stream, a first-pass processing will produce detector conditions and calibration constants of sufficient quality needed for bulk calibration of all streams. Finally, full events are calibrated and reconstructed in as much detail as possible during offline processing to be delivered for analyses.

The full offline reconstruction starts from detector-specific objects, such as calorimeter clusters and charged particle tracks, to identify all physics objects in the events. The main physics stream is fully reconstructed to produce Event Summary Data (ESD), which contains detailed reconstruction output, such as calorimeter cells and track hits. Most physics analyses do not demand this level of detail since reconstructed physics objects suffice, hence, unless otherwise requested the ESD typically is deleted as soon as T0 subsequently converts it into Analysis Object Data (AOD), which is widely recognized as the primary format within ATLAS. These data are forwarded to filters and derivation frameworks at T0 to produce derived data in the form of Derived RAW data, Derived Event Summary Data (DESD) or Derived Analysis Object Data (DAOD). Both the original and derived data are saved to tape at T0; the derived data are also distributed to Tier-1 data centers.

3.4 Monte Carlo Simulation

The ATLAS simulation [167] is a Monte-Carlo-technique-based, multicomponent production chain, from generators that produce partial or complete events, to packages that simulate the response of subdetectors and triggers, and to algorithms that reconstruct analyzable objects. It is seamlessly integrated into the ATLAS software framework, Athena. Its output, in a format identical to that of the true detector, makes it ready for large-scale

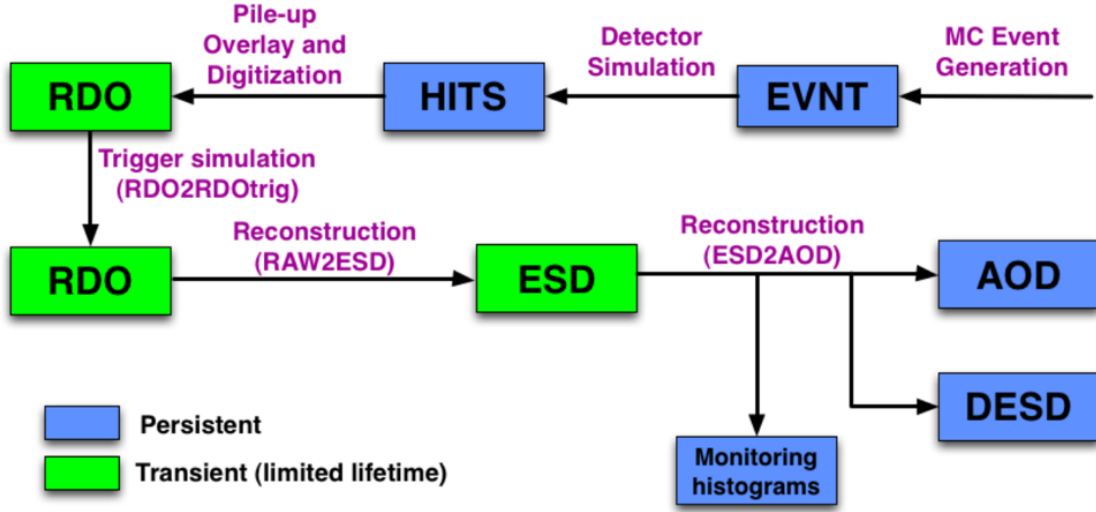


Figure 3.12: ATLAS Monte Carlo Simulation Data Flow [168].

distribution on the WLCG. Figure 3.12 highlights the main components and output formats.

3.4.1 Event Generation

Hadronic collisions typically are not modelled by one single event generator. The simulation of hard scatters (hardest parton-parton interactions that produce signal such as HIPs) models kinematics, initial-state radiation, final-state radiation and hadronization of partons. Moreover, underlying events, that is, interactions of other partons that do not take part in the hard scatters, need to be modelled adequately. Event generators tend to focus on different event components, and thus can be categorized accordingly. General-purpose generators produce complete events from initial states such as proton-proton collisions. They are either used standalone or in conjunction with specialized generators that target specific components of the events.

MADGRAPH5_aMC@NLO [147] is a specialized generator that calculates matrix elements for decays and $2 \rightarrow n$ scatters. Given a user-defined process in terms of initial- and final-

state particles, it generates Feynman diagrams for the process, and prepares the matrix element for cross section or decay-width calculations. A parton distribution function (PDF) must be provided as an input to describe the energy spectrum of the initial partons. In ATLAS, this is supported by a repository of PDFs via the Les Houches Accord PDF (LHAPDF) interface [149]. Kinematic properties of the initial- and final-state particles are stored in the Les Houches Events (LHE) files. PYTHIA8 [169], as a general-purpose generator, showers and hadronizes partons saved in the LHE files. The modeling of underlying events is controlled by adjustable parameters from fits to real data, known as a tune. The soft inelastic proton-proton interactions or pileup events are generated with PYTHIA8, according to a parton distribution function and a tune describing the pileup profile of a given data-taking period.

A particle-gun event generator produces events with single particles at specified energies and directions. It proves useful for detector calibration and reconstruction efficiency studies. Additionally, it can be exploited to loosely estimate efficiencies and limits for models with similar signatures to a specific model. In addition to not being based on a valid physics model, a particle-gun generator does not produce underlying events, proton remnants remaining after collisions, or other primary interactions between the colliding protons.

3.4.2 Detector Simulation

The detector simulation software processes truth information of particles from the event generators through the GEANT4 simulation toolkit [170], which itself is embedded in the Athena framework with access to databases of detector geometry, i.e., geometrical layout and atomic composition, as well as conditions such as calibration, dead channels, etc. The numerous GEANT4 models for particle interactions are organized by *physics lists*, which are distributed as source code of GEANT4. Users have different modelling descriptions at their disposal by choosing different *physics lists*. An implication is that when any of the

models is modified, a new *physics list* will need to be created. Some physics processes, for instance, ionization and bremsstrahlung, have large cross sections at low energies, therefore, production cuts are implemented in GEANT4 so that computational resources are not spent on the propagation of low-energy particles. As incident particles step through the detector with trajectories simulated according to their equations of motion, their energy deposits (HITS) in sensitive detector volumes are computed according to their rate of energy loss and subtracted from their initial energies. The hard-scatter interactions and the pileup interactions are separately simulated and stored.

3.4.3 Digitization and Reconstruction

The purpose of digitization is to convert the HITS into digital detector responses (digits) to be used as input to the reconstruction algorithms. Signal HITS from the hard process are read in simultaneously with pileup HITS, so the energy deposits of these particles are combined or overlaid for each detector element, and then passed to the digitization. The readout of most subdetectors are affected by interactions from neighbouring bunch crossings, thus, it is a necessity to read in the out-of-time pileup events in addition to the in-time pileup events for any given hard-scatter event. Each of the neighbouring bunch crossings is assigned the same value of μ (see definition in Section 3.1) as the current bunch crossing, although in data there are variations between adjacent bunches. The number of pileup events for each bunch crossing is randomly drawn from a Poisson distribution with a mean of μ , while the events themselves are sampled from the pileup HITS file based on their cross sections [171]. The detector response is subsequently emulated for HITS from all contributing bunch crossings set by the longest readout time window. The resulting digits are sent to software emulation of detector readout electronics, whose output is recorded as Raw Data Output (RDO). As the final output of the digitization step, the RDO has the

same format as that of real collision data and thus enables the execution of the 2-tier trigger selection through emulation as well as reconstruction algorithms described in the preceding subsection, through the Athena framework.

3.5 Data Storage and Analysis

Although collision data in its various formats undergo an initial processing at CERN T0, their storage and analysis are mainly managed by the three lower tiers of the WLCG, as is the handling of the MC simulation data.

As described in Section 3.3.2, T1 nodes around the globe receive RAW and reconstructed data from CERN T0. They reprocess the RAW data into (D)ESD or (D)AOD formats using updated calibration constants or reconstruction algorithms. This is CPU-intensive and thus is typically performed during dedicated yearly campaigns. In addition to the data reprocessing, the MC simulation also takes place at T1, and in some cases at Tier-2 (T2).

The last two tiers, composed of T2 institutional computing centres and Tier-3 (T3) local clusters, are capable of rewriting disk data on a regular basis, therefore, they are designated for analysis purposes. Except for the RAW data that are restricted to T0 and T1 sites, the other formats can be distributed to the lower tiers, where a wide range of analysis jobs are launched. One such job is to convert the (D)ESD and (D)AOD into an ntuple to store data in vectors rather than as objects. This is done to further reduce file size through selected triggers, selection cuts, etc., for the ease of data analysis.

Chapter 4

Simulation of Highly Ionizing Particles

Highly ionizing particles (HIPs) must leave some signature in a measurable way to be detected. Based on their interaction with matter, a full simulation is exploited to study their signature. The aspects germane to the HIP simulation in ATLAS are described in this chapter.

4.1 Interaction with Matter

The detection of HIPs relies on their electromagnetic interaction with matter. Direct interaction with nuclei is shielded by atomic electrons and thus only takes place when the HIPs have sufficiently high momenta. Therefore, as other charged particles do, the HIPs interact primarily with the atomic electrons, through which scattering and energy loss occur. If a charged particle is massive, its collision with atomic electrons, as described by Rutherford scattering, will not substantially alter its trajectory or contribute to its energy loss. However, if the particle has adequate kinetic energy, it can ionize atoms and lose that energy until

it comes to a halt. When the particle accelerates in the atomic electric field, it also loses its kinetic energy through radiation of electromagnetic waves. The emission of photons, accompanied by deceleration of the particle, is known as bremsstrahlung. For HIPs within the kinematic range accessible at the LHC, the ionization energy loss predominates over the bremsstrahlung radiation.

The calculation of HIP interactions with matter presents formidable challenges. The long-range nature of the electromagnetic force necessitates the treatment of HIP interaction with targets as two-body collisions with distance dependence. Also, the application of quantum mechanics is crucial in deriving solutions to the interaction.

4.1.1 Rutherford Scattering

Rutherford scattering [172] is the process of elastic scattering between non-relativistic charged particles by the Coulomb interaction. In the event of HIP scattering off an atomic electron, Rutherford scattering not only suggests the β -dependent magnetic coupling to be mentioned, with β being the HIP speed $v = |\vec{v}|$ normalized by the speed of light c , it also helps clarify the subsequent calculation of HIP energy loss.

The two interacting particles separated by a distance $r = |\vec{r}|$ experience a Coulomb force that gives rise to an electric potential

$$V(r) = \begin{cases} -\frac{ze^2}{r} & \text{for a HECO with charge } ze, \\ -\frac{eg}{c} |\vec{v} \times \frac{\vec{r}}{r^2}| & \text{for a monopole with charge } g. \end{cases} \quad (4.1)$$

The HIP interaction with an atomic electron can be analyzed in polar coordinates with the origin at the electron's position (see Figure 4.1). Given a specific form of $V(r)$, one can determine the HIP coordinate $\vec{r}$, and therefore the scattering angle θ , as a function of the

impact parameter b . In the case of HECOs, b and θ are related to each other by

$$b = \frac{ze^2}{2E} \cot \frac{\theta}{2}, \quad (4.2)$$

where $E = \frac{1}{2}mv^2$ is the initial energy of the incident particle with mass m .

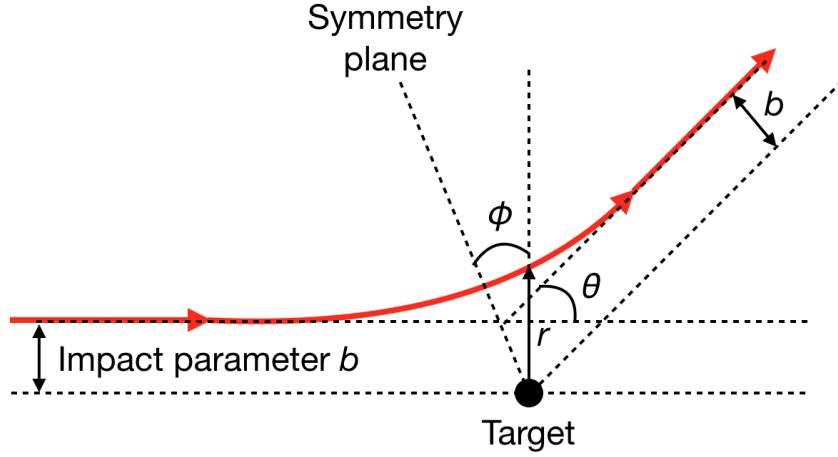


Figure 4.1: Depiction of Rutherford scattering of a charged particle in motion (red line) with a stationary charged particle denoted as target. The azimuthal angle ϕ fixes the position of the symmetry plane. The polar angle θ is the angle between the horizontal axis and the charged particle in motion.

The differential cross section of a HIP scattered by the Coulomb field of an electron is inseparably linked to the specific relationship between b and θ . As the scattering source, the electron has an effective area of $2\pi b db$ at impact parameter b , for the HIP to be scattered by the angle θ into a solid angle $d\Omega = \sin \theta d\theta d\phi$. Since the scattering has no azimuthal dependence, the differential cross section only depends on the angle θ and not on ϕ . Integrating over ϕ , one obtains the cross section for scattering through an angle θ :

$$- \int_0^{2\pi} d\phi \frac{d\sigma}{d\Omega}(\theta) \cdot \sin \theta d\theta = 2\pi b db, \quad (4.3)$$

where the negative sign signifies the anti-correlation between b and θ . Rearranging this

equation gives the differential cross section

$$\frac{d\sigma}{d\Omega}(\theta) = -\frac{b}{\sin\theta} \frac{db}{d\theta}. \quad (4.4)$$

Based on Equation 4.2, the HECO differential cross section is

$$\frac{d\sigma_{\text{HECO}}}{d\Omega} = \left(\frac{ze^2}{4E}\right)^2 \frac{1}{\sin^4(\theta/2)}. \quad (4.5)$$

It can be readily seen that most HECOs will have very little angular deviation θ . For small scattering angle $\theta \ll 1^\circ$ and singular electric charge $z = 1$, the electron differential cross section is

$$\lim_{\theta \ll 1} \frac{d\sigma_{\text{electron}}}{d\Omega} = \left(\frac{e^2}{2mv^2}\right)^2 \frac{1}{(\theta/2)^4}. \quad (4.6)$$

The electron-monopole scattering exhibits a different θ dependence in general. However, in the limit of small scattering angles [142], the monopole potential given by Equation 4.1 reduces to $V(r) = -\frac{egv}{cr}$, which leads to the differential cross section of

$$\lim_{\theta \ll 1} \frac{d\sigma_{\text{monopole}}}{d\Omega} = \left(\frac{eg}{2mvc}\right)^2 \frac{1}{(\theta/2)^4}. \quad (4.7)$$

It is apparent that Equation 4.7 becomes identical to Equation 4.6 if $\frac{e}{v}$ is used in place of $\frac{g}{c}$.

The differential cross section of the Rutherford scattering can be expressed for, instead of a solid angle, a given amount of energy transfer. For HECOs, this is

$$\frac{d\sigma}{dT} = \frac{2\pi(ze^2)^2}{m_e v^2} \frac{1}{T^2}, \quad (4.8)$$

which is interpreted as the probability for a HECO of charge z and velocity v to lose kinetic energy T when colliding with a spinless electron.

It should be noted that the differential cross section of Rutherford scattering is subject to

four major assumptions. Foremost, the first-order plane-wave Born approximation is applied to the incident and scattered HIP; this approximation can be made as long as the incident energy is large and the scattering potential is weak. Second, the target electron is so heavy that it does not recoil as the HIP scatters; this is satisfied in elastic scattering if the incident HIP transfers sufficiently small energy in comparison to the rest energy of the target electron. In addition, the incident HIP and target electron are assumed to be spin-0 point particles. These restrictions certainly invite justification or removal, as will be discussed shortly.

4.1.2 Energy Loss

The consequence of a large number of energy transfers, each described by Equation 4.8 for instance, is the continuous deceleration of the HIP. The process of deceleration can be characterized schematically by the amount of kinetic energy lost, $-dT$, per unit path length, dx , in a material of electron number density n :

$$-\frac{dT}{dx} = n \int_{T_{\min}}^{T_{\max}} T \frac{d\sigma}{dT} dT, \quad (4.9)$$

where the integral, bounded by an upper limit of energy transfer $T_{\max}$ and a lower limit of $T_{\min}$, sums energy loss per scattering with a weight of its differential cross section $\frac{d\sigma}{dT}$. In practice, the energy loss is not measured in length units dx but in terms of $d(\rho x)$, where ρ is the material density. The energy loss, also referred to as the stopping power, is then

$$-\frac{dT}{d(\rho x)} = -\frac{1}{\rho} \frac{dT}{dx} \equiv -\frac{dE}{dx}. \quad (4.10)$$

Of practical concern is the restricted validity of this formulation. The electron spin should be considered. Its effect arises from the incident particle's magnetic moment interacting with the target electron's magnetic field, which vanishes in the rest frame of the target and

is maximal in the rest frame of the incident particle. Absent in the classical Rutherford approach, the electron spin is accounted for in Mott scattering [173]. The range of energy transfer must be defined. In the limit of large energy transfer, set by $T_{\max}$ in Equation 4.8, the electrons are assumed to be free, as their binding energy is much smaller than the energy transfer. The upper bound can be computed using the kinematics of free electrons in hard collisions. However, the electrons are no longer free in the event of small energy transfers, when their binding energy dominates. In the regime of soft collisions, set by $T_{\min}$, the solution of Equation 4.8 must involve a description of the atomic structure, approximated by the atomic charge number Z and mass number A .

Electromagnetic Ionization

As a matter of fact, the HIP ionization is treated in two kinematic regimes. In soft collisions with energy transfers on the order of $T_{\min}$, excitation of the bound electrons leads to a Z -dependent mean excitation potential. Since these electrons are not stationary, their kinematics modify the energy loss through a shell correction that is attributed to each electron shell. On the other hand, in the regime of high energy transfers, the HIP electric fields undergo considerable Lorentz contractions, and thus polarize the atoms, even more so in denser materials. Handled by a density correction, the outcome is that the HIPs experience less energy loss in collisions with electrons shielded from the full electric fields of the HIPs.

For HECOs, these effects are taken into account by the Bethe-Bloch formula of energy loss [174, 175],

$$-\frac{dE}{dx} = Kz^2 \frac{Z}{A} \frac{1}{\beta^2} \left[\ln \frac{2m_e c^2 \beta^2 \gamma^2}{I} - \beta^2 - \frac{\delta(\beta\gamma)}{2} \right], \quad (4.11)$$

with $K = 4\pi N_A r_e^2 m_e c^2 = 0.307 \text{ MeV g}^{-1} \text{ cm}^2$ and density correction $\frac{\delta(\beta\gamma)}{2} = \ln \frac{\hbar\omega_p}{I} + \ln \beta\gamma - \frac{1}{2}$, where $\omega_p = 28.16 \sqrt{\rho Z/A}$. The HECO electric field polarizes the atoms longitudinally such that the HECOs lose less energy when colliding with distant electrons than with close

electrons. The shell correction is dropped in Equation 4.11, as it is small in general.

Peering beyond HECOs into monopoles, a different stopping power is called for. The Ahlen-Bloch formula [40], integrating close interactions modelled by Kazama, Yang and Goldhaber [176] and distant interactions modelled using a dipole approximation, describes energy loss of monopoles with velocities $\beta \geq 0.2$ and Lorentz factor $\gamma \leq 100$, a kinematic regime where the spin effect is considered negligible.

The energy loss of magnetic monopoles acquires a charge dependence of g^2 , as has been noted during the computation of differential scattering cross sections in the small-angle limit (see Equation 4.7). Although the mean excitation potential and shell correction still apply, the density correction is less significant, as there is no longitudinal interaction with atomic electrons via the transverse electric field produced by a moving magnetic monopole. The monopole energy loss is derived as

$$-\frac{dE}{dx} = K \frac{Z}{A} g^2 \left[\ln \frac{2m_e c^2 \beta^2 \gamma^2}{I} + \frac{k(g)}{2} - \frac{1}{2} - \frac{\delta}{2} - B(|g|) \right], \quad (4.12)$$

with the Kazama correction of the QED radiative effect

$$k(|g|) = \begin{cases} 0.406 & \text{for } |g| = 1g_D \\ 0.346 & \text{for } |g| = 2g_D \end{cases} \quad (4.13)$$

that entails the emission of photons as the electromagnetic field around the monopole rearranges, and the Bloch correction

$$B(|g|) = \begin{cases} 0.248 & \text{for } |g| = 1g_D \\ 0.672 & \text{for } |g| = 2g_D \end{cases} \quad (4.14)$$

that encodes the different cross-sectional areas of the incident HIP beam relative to the

target.

Comparing monopoles and HECOs, one sees in Figure 4.2 that their energy losses diverge as their velocity decreases. First, HECOs deposit more of their kinetic energy at lower velocities, as their energy loss varies as a function of β^{-2} . Consequently, the energy loss acquires an implicit dependence on particle mass due to the relationship between velocity and mass, i.e., $1/\beta^2 = m^2\gamma^2/p^2$. That is, energy loss incurred at low velocity ($\gamma \approx 1$) will not be the same for HIPs of the same momentum p but different mass m . This is of particular consequence for HECOs, whose dE/dx values scale quadratically as mass and energy, whereas for magnetic monopoles, the leading-order contribution is logarithmic and thus is weaker. In contrast to the monopole case, the HECO energy loss is much more susceptible to velocity increase, as its β^{-2} -dependent drop is suppressed by an onset of a logarithmic rise due to relativistic effects. At sufficiently high $\beta\gamma$, the ionization energy loss for both monopoles and HECOs, dominated by the logarithmic rise, becomes saturated. This saturation, as a consequence of large energy transfer in a cascade of ionization events spanning a long distance, marks the transition to substantial production of δ -rays in substantial quantities – a phenomenon that gains more prominence as the HIPs become relativistic.

The energy loss accumulates uncertainties from various sources listed in Table 4.1. A total uncertainty of 3% accrues for both monopoles and HECOs.

	Monopole	HECO
Mean excitation potential	1%	1%
Shell correction	1%	1%
Density correction	2%	< 1%
Kazama correction	-	< 1%
Bloch correction	-	1%
Total	$\leq 3\%$	$\leq 3\%$

Table 4.1: Summary of HIP dE/dx uncertainties [40, 41].

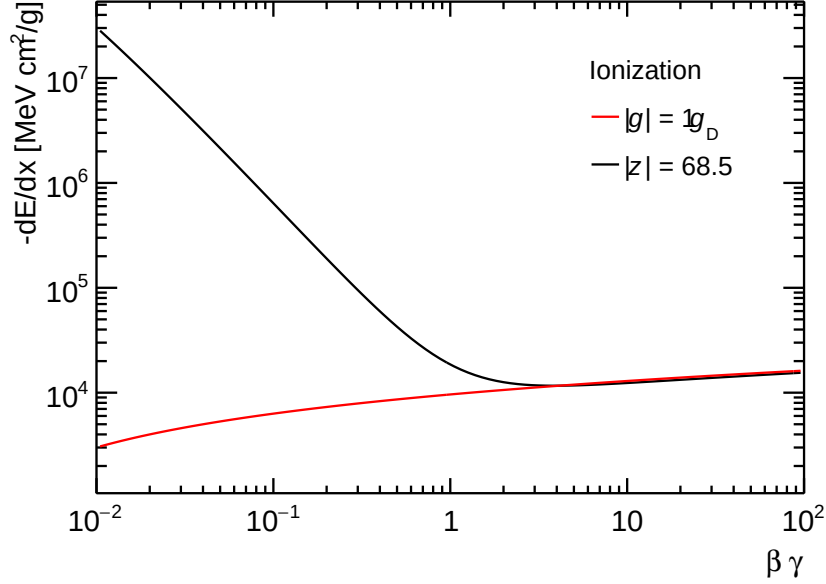


Figure 4.2: HIP ionization energy loss as a function of $\beta\gamma$ for monopoles with charge $|g| = 1g_D$ and HECOs with charge $|z| = 68.5$. Silicon is chosen for the numerical computation, although the material only has a minor effect on the energy loss.

Bremsstrahlung

In addition to ionization energy loss, fast charged particles lose energy by photon emission as they are accelerated in the Coulomb field of the nucleus. The bremsstrahlung energy loss for high-energy particles depends on their charge z , mass m and energy E :

$$-\frac{dE}{dx} \propto \frac{z^2}{m^2} E. \quad (4.15)$$

It is proportional to the energy of the particle, while the ionization energy loss is proportional to the logarithm of the energy in the relativistic regime. Therefore, bremsstrahlung dominates in the ultrarelativistic regime.

Despite the m^{-2} dependence, bremsstrahlung can be noticeable for massive particles if they have sufficiently high energies. This can be explained with the concept of critical energy, at which energy loss due to ionization equals that due to bremsstrahlung. The critical energy

E_c^μ of muons in iron is 890 GeV [177]. For HECOs with mass 200 GeV and charge $|z| = 100$, the critical energy in the same medium is given by

$$E_c^{\text{HECO}} = E_c^\mu \left(\frac{m_{\text{HECO}}}{m_\mu} \right)^2 \left(\frac{z_\mu}{z_{\text{HECO}}} \right)^2 = 3.2 \times 10^4 \text{ TeV}, \quad (4.16)$$

indicating in effect a negligible energy loss due to bremsstrahlung for the HIPs probed at the LHC.

Energy-Range Relations

Having computed the HIP energy loss, one can estimate the HIP range, namely the distance the HIP is expected to travel before it stops. Given an initial energy E_0 , the range is

$$R(E_0) = \int_{E_0}^m \frac{dE}{dE/dx}, \quad (4.17)$$

where m is the HIP mass.

The HIP range inherits the dE/dx mass dependence. Building on the discussion thus far, one can see that with different masses, low-energy HIPs ($\gamma \approx 1$), especially HECOs, can have radically different ranges even if they have the same kinetic energy. High-energy HIPs ($\beta \approx 1$) of the same charge, on the other hand, have approximately the same range, which is proportional to their energy, regardless of their masses.

4.2 Simulated Detector Response

The HIP ionization properties profoundly impact the detector response, as illustrated in Figure 4.3 for a magnetic monopole. A sequence of TRT high-threshold hits in red not only measure the position but also reflect the energy loss of the particle to the atomic electrons in the detector material. Equally evident is a concentrated energy in the Electromagnetic

Calorimeter, depicted by a green wedge. To unravel the event topology, GEANT4 simulation packages are exploited to model particle propagation and energy deposition using specialized functions and tools.

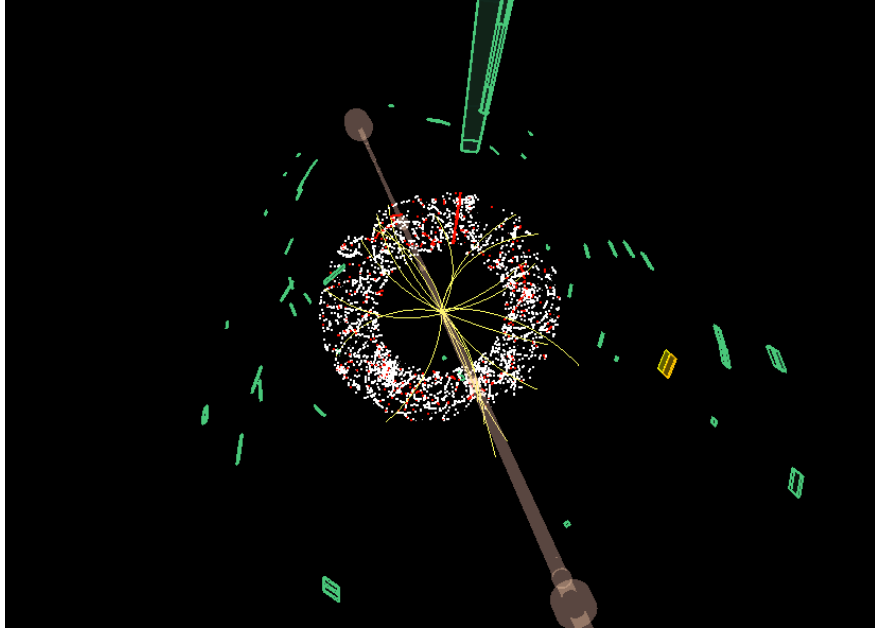


Figure 4.3: Event display of a $1g_D$ $m = 1500$ GeV monopole in a particle-gun event [74].

4.2.1 Trajectory

The electrically neutral monopoles and the electrically charged HECOs will experience different forces while propagating within the ATLAS magnetic field. A monopole with magnetic charge g is accelerated by the solenoidal magnetic field, $\vec{B} = B\hat{z}$:

$$\frac{d\vec{p}}{dt} = g\vec{B}, \quad (4.18)$$

where $\frac{d\vec{p}}{dt}$ is the time evolution of the monopole momentum $\vec{p}$. With the ATLAS coordinate system (see Section 3.2), Equation 4.18 indicates that the monopole trajectory is straight in the $r - \phi$ plane; the bending of the monopole happens in the $r - z$ plane as the magnetic field

changes the monopole pseudorapidity along its trajectory. On the other hand, the HECO's equation of motion reads

$$\frac{d\vec{p}}{dt} = q\vec{v} \times \vec{B}, \quad (4.19)$$

for a HECO with electric charge q . In the absence of acceleration by the magnetic field, the HECO trajectory is straight in the $r - z$ plane and curves in the $r - \phi$ plane, given by the time evolution of the HECO momentum in Equation 4.19. In addition, the HECO experiences stronger bending than Standard-Model particles of the same velocity, since the bending scales as the electric charge of the particle, as per Equation 4.19.

Equation 4.18 and Equation 4.19 are applied in the ATLAS simulation to propagate monopoles and HECOs, respectively. Invoked in the `G4mplEquationSetup` class in the Athena package `Simulation/G4Extensions/Monopole`, the GEANT4 `ClassicalRK4` stepper employs Equation 4.18, while the default `NystromRK4` stepper is based upon Equation 4.19 for electrically charged particles [178].

The straight monopole trajectory in the $r - \phi$ plane dictates that the difference between the electromagnetic cluster ϕ coordinate and the monopole truth ϕ should be small. This difference is shown as a function of monopole truth kinetic energy in Figure 4.4 for particle-gun monopoles of mass 200 GeV and for magnetic charges $g = \pm 1g_D$ separately. With a fixed truth kinetic energy, lighter HIPs are more likely to deposit sufficient energies and be reconstructed as clusters in the Electromagnetic Calorimeter. Thus, the mass of 200 GeV, the lowest considered in the analysis detailed in the subsequent chapter, was chosen to ensure that the cluster η is well defined. On the other hand, the bending monopole trajectory in the $r - z$ plane manifests as the difference between the electromagnetic cluster η coordinate and the monopole truth η . This can be seen in Figure 4.5 for the monopoles with truth $|\eta| < 0.3$. This relatively low truth $|\eta|$ range was chosen to accentuate the change in η as the monopoles propagate, given the solenoidal magnetic field being directed along $|\eta| = \infty$.

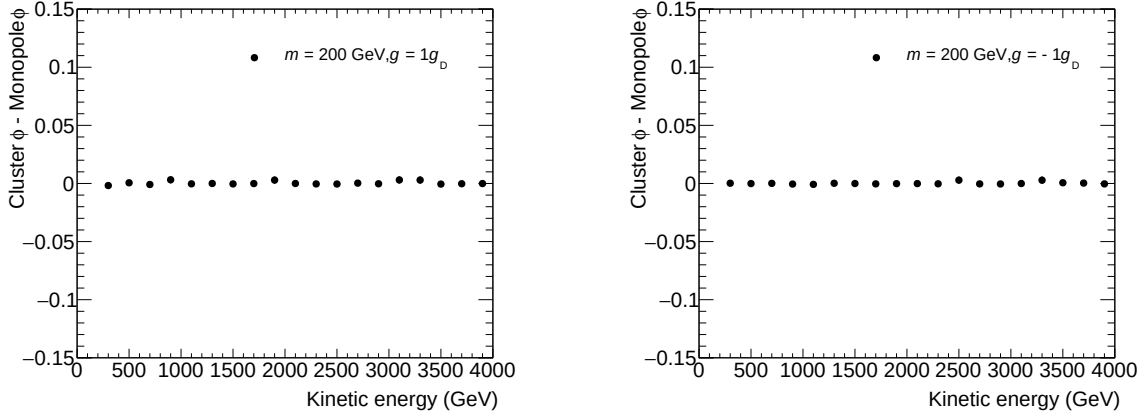


Figure 4.4: The difference between the EM cluster azimuthal coordinate ϕ and the truth ϕ as a function of truth kinetic energy for $m = 200 \text{ GeV}$ monopoles with magnetic charge $g = 1g_D$ (left) and with magnetic charge $g = -1g_D$ (right).

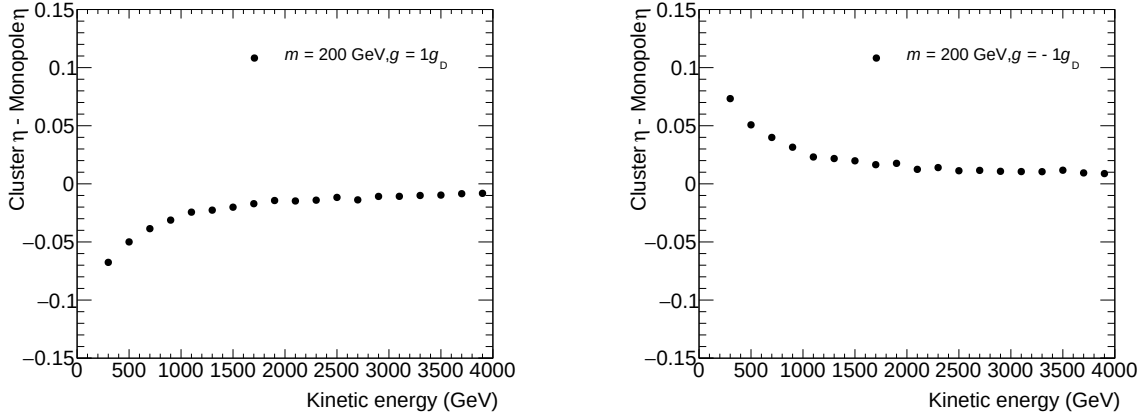


Figure 4.5: The difference between the EM cluster pseudorapidity η and the truth η as a function of truth kinetic energy for $m = 200 \text{ GeV}$ monopoles with magnetic charge $g = 1g_D$ (left) and with magnetic charge $g = -1g_D$ (right).

The HECO trajectories are characterized in Figures 4.6 – 4.7, as the spatial differences between the electromagnetic clusters and the truth particles. While the η coordinate does not change, the ϕ coordinate varies quite significantly.

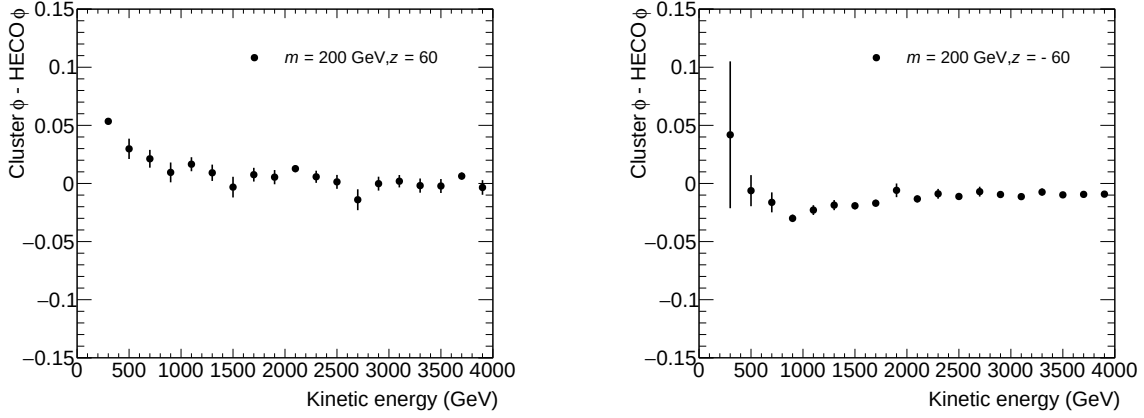


Figure 4.6: The difference between the EM cluster azimuthal coordinate ϕ and the truth ϕ as a function of truth kinetic energy for $m = 200$ GeV HECOs with electric charge $z = 60$ (left) and with electric charge $z = -60$ (right). The same range is used as in Figure 4.4 for the purpose of comparison.

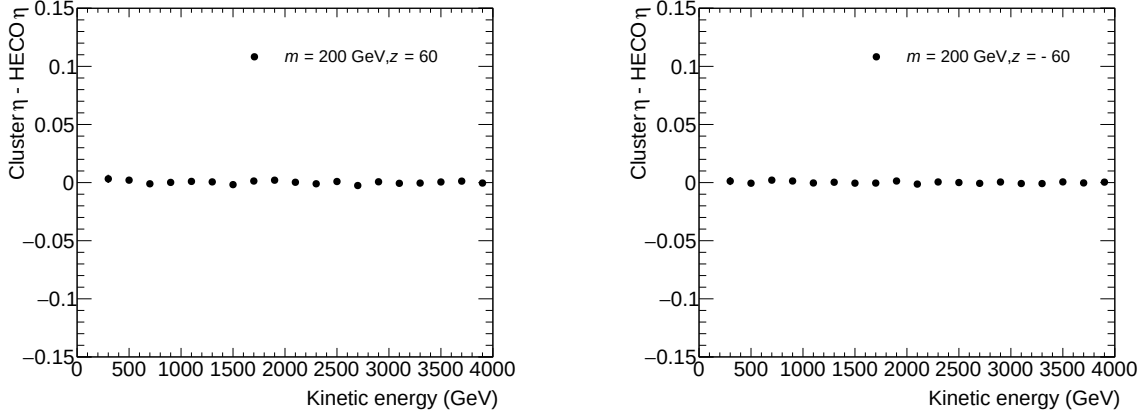


Figure 4.7: The difference between the EM cluster pseudorapidity η and the truth η as a function of truth kinetic energy for $m = 200$ GeV HECOs with electric charge $z = 60$ (left) and with electric charge $z = -60$ (right). The same range is used as in Figure 4.5 for the purpose of comparison.

4.2.2 Energy Deposition and δ -Ray Production

When a HIP travels through the ATLAS Detector, it deposits energy via ionization and δ -ray production, both of which are simulated by dedicated GEANT4 classes in the package Simulation/G4Extensions/Monopole. The class G4mplAtlasIonisation is intended for the simulation of energy loss by magnetic monopoles to atomic electrons in the de-

tector material, while the class `G4hIonisation` is utilized for HECOs. Implemented in the `G4mplAtlasIonisationWithDeltaModel` class, the production of δ -rays is sampled from a probability distribution. Once a δ -ray is created, its energy is subtracted from the energy of the primary HIP, resulting in a decrease of the HIP dE/dx . On the contrary, if the δ -ray production is suppressed, the short-distance primary ionization is enhanced, manifesting in an increase of the HIP dE/dx .

4.2.3 Inner Detector Hits

The Inner Detector records signal hits if a charged particle deposits enough energy to overcome some noise-discriminating threshold. As HIPs deposit a large amount of energy that exceeds the saturation limit of the Pixel detector and the Insertable B-Layer (IBL), they result in overflow bits in these detectors. Since the Pixel detector does not register the overflow and sets the bits to zero, the information on the deposited energy is lost in the detector readout. Despite the IBL records multibit digits that include the overflow bits, this information hasn't proved to outperform the TRT hit information to be discussed shortly, hence, it is not exploited in this analysis. The HIPs are expected to leave large clusters in the SemiConductor Tracker; the discriminating power of these clusters has not been investigated and is worth further examination in the future. Nevertheless, the HIP energy depositions are distinguishable in the Transition Radiation Tracker (TRT) with its high readout thresholds. A HIP will cross 35 TRT straws on average, giving rise to 35 high-threshold (HT) hits by its primary ionization and additional HT hits by the simultaneous formation of δ -rays. Although the production of δ -rays has a larger cross section in the straws than in the radiator material due to the higher density of the straw gas than that of the radiator material, the δ -rays produced in radiator material propagate to nearby straws and thus give rise to additional HT hits. The HIP ionization in the TRT is modelled using

the GEANT4 simulation toolkit, although for Standard-Model charged particles the GEANT4 ionization is overwritten in the TRT digitization with the Photo-Absorption Ionization (PAI) model [179] to distribute energy deposits along particle trajectories in the straws.

The distribution of δ -rays lends itself to constructing a region in the TRT with HT hits from the HIP traversal. This distribution can be predicted by the δ -rays' range and radius of curvature. The bending of a δ -ray, represented by the radius of curvature r , in a magnetic field B , is related to its momentum p and charge e , via $p = eBr$, which reduces to p [MeV] = $0.3 B$ [T] r [mm] in the 2-T axial magnetic field. That is, a typical δ -ray with 3-MeV momentum will have a curvature of about 5 mm. Since this δ -ray has a longer range than 5 mm, its radius of curvature determines its relative distance to the HIP. Therefore, a host of δ -rays produced by the HIP ionization are contained in a ± 5 mm region around the HIP trajectory, in the form of TRT HT hits.

4.2.4 Liquid Argon Charge Recombination

When an incident particle ionizes the argon atoms in the Electromagnetic Calorimeter, the current detected is the only attainable information about the energy deposited by that particle. In order to accurately measure that energy, the electron-ion pairs must be separated by an electric field as soon as they are formed, thereby keeping them from recombining into atoms. The recombination effect also depends on the material density and the incident particle's dE/dx , both of which intensify the formation of electron-ion pairs. Although ideally the true energy deposition should be reflected by the actual number of electrons liberated by ionization, in practice, one can only measure the current from a reduced number of electrons as a result of charge recombination. This measurement gives rise to a visible energy smaller than the true energy one wishes to measure.

HIPs ionize the liquid argon atoms and produce a great many electrons and ions sur-

rounded by δ -rays in a narrow region. The reduction of detectable electrons and the underestimation of energy loss to the atomic electrons incurred by the recombination effect must be accounted for.

Theory and Experiment

Theoretical treatment of this phenomenon draws upon assumptions on the initial conditions of the electrons and positive ions, in order to calculate their time and spatial evolution under an external electric field. The effectively straight track of a heavy charged particle constitutes a core, surrounded by a number of electron-ion pairs. As the electrons and positive ions drift towards their corresponding electrodes, their density distributions smear out, and their number densities diminish upon recombination. There are many attempts to model the spatial distributions of these ions and their variation in time [180–182], for example, the one developed by Jaffé in 1913 [180]. His approach imposes a cylindrical volume around the particle trajectory to encompass the ionization electrons and positive ions, which are modelled by Gaussian distributions around the trajectory of the incident particle with the same initial distribution. The fraction of electrons or positive ions that remain visible at infinite time is shown to be

$$R = \frac{1}{1 + k \frac{dE/dx}{E_D}}, \quad (4.20)$$

where k is a material-dependent constant. Now referred to as Birks' Law [183, 184], initially it was used to model the scintillation effect of ionizing particles. Compared to the other theoretical models [181, 182], the Jaffé model has a relatively simple functional form, and it has important applications in describing ionization interactions on account of its dE/dx dependence.

This indeed is attested to by the ICARUS Experiment when investigating the response of their LAr Time Projection Chamber in a drift electric field range $0.1 < E_D < 1.0$ kV/cm to

cosmic-ray muons and protons in an ionization density range $2.1 < dE/dx < 42 \text{ MeV/cm}$ [185].

A fit to Birks' Law of the functional form

$$R = \frac{A}{1 + k \frac{dE/dx}{E_D}} \quad (4.21)$$

yields a normalization constant $A = 0.800 \pm 0.003$ and the Birks' constant $k = 0.0486 \pm 0.0006 \text{ kV/cm} \frac{\text{g/cm}^2}{\text{MeV}}$. The interpretation of this ratio is complicated by the presence of δ -rays. In the actual detector, the produced δ -rays, as part of the true energy deposition E_0 , will contribute to the energy deposition measured by the detector, or visible energy deposition E_{vis} , as long as they do not recombine with positive ions. Since the δ -rays are indistinguishable from primary ionization, the ratio in data can be trivially defined:

$$R_{\text{Exp}} = \frac{E_{\text{vis}}}{E_0}. \quad (4.22)$$

In simulation, however, the δ -rays are not treated on an equal footing as the primary ionization. A certain energy threshold, realized via a cut in range (see Equation 4.17), is set to forbid explicit simulation of low-energy δ -rays. Thus, the visible energy deposition in MC is composed of $E_{\text{vis,ion}}$ from the primary ionization and $E_{\text{vis,delta}}$ from the δ -rays, that are separately sampled from their corresponding probability distributions, leading to a different expression of

$$R_{\text{MC}} = \frac{E_{\text{vis,ion}} + E_{\text{vis,delta}}}{E_0}. \quad (4.23)$$

Variation of the energy threshold will shift the relative proportions of the two contributions.

ATLAS Implementation

Relating the visible deposited energy E_{vis} to the true deposited energy E_0 for particles with energy loss dE/dx in the LAr calorimeter of material density ρ and drift electric field E_D ,

ATLAS implements Birks' Law in the class LArG4BirksLaw in the package LArCalorimeter/LArG4/LArG4Code, as

$$E_{\text{vis}} = E_0 \frac{1 + A'k/E_D}{1 + k \frac{dE/dx}{\rho E_D}}. \quad (4.24)$$

The Birks' constant k is set to the ICARUS measurement. The constant A' is set to normalize the fraction of visible energy E_{vis}/E_0 to that of a minimally ionizing particle that satisfies $dE/dx/\rho = 1.51$. Assuming $E_{\text{vis}} = E_0$ for such a particle, the normalization parameter $A' = 1.51$.

The ATLAS implementation in Equation 4.24 is not directly applied to HIPs, whose high dE/dx is beyond the range probed by ICARUS and needs to be treated with a different functional form than Equation 4.24. As a result, it is modified to take the functional form

$$E_{\text{vis}} = k_{\text{hip}} \left(E_0 \frac{1 + A'k/E_D}{1 + k \frac{dE/dx}{\rho E_D}} \right), \quad (4.25)$$

where a dE/dx -dependent HIP correction k_{hip} is derived from a comparison of MC to published heavy-ion data, by tuning the MC recombination factor R_{MC} to the data recombination factor R_{data} (see [186] and references therein).

The computation of R_{MC} involves separating the energy depositions from HIP ionization and δ -rays have drastically different dE/dx values. Since the correction k_{hip} scales the HIP visible energy $E_{\text{vis,hip}}$, the MC recombination factor becomes

$$R_{\text{MC}} = \frac{k_{\text{hip}}E_{\text{vis,hip}} + E_{\text{vis,delta}}}{E_0}. \quad (4.26)$$

This allows a rearrangement of

$$\frac{R_{\text{MC}}}{R_{\text{Exp}}} = \frac{k_{\text{hip}}E_{\text{vis,hip}} + E_{\text{vis,delta}}}{E_0 R_{\text{Exp}}} = 1, \quad (4.27)$$

from which the HIP correction is extracted,

$$k_{\text{hip}} = \frac{R_{\text{Exp}} E_0 - E_{\text{vis,delta}}}{E_{\text{vis,hip}}}. \quad (4.28)$$

This correction is calculated for the different heavy-ion beams in the data, using the experimental measurement R_{Exp} and E_0 , $E_{\text{vis,delta}}$, and $E_{\text{vis,hip}}$, computed from ATLAS simulation. Any change to the relevant simulation parameters, i.e., parameters in Equation 4.24 and the electron range cut in the LAr Electromagnetic Calorimeter, would require the appropriate HIP correction to be rederived.

A linear function of dE/dx is fitted to the k_{hip} values computed for $969 < dE/dx < 12,000$ MeV/cm, within which the heavy-ion data is collected. For dE/dx values below 969 MeV/cm and above 12,000 MeV/cm, k_{hip} is set to its values at $dE/dx = 969$ MeV/cm and $dE/dx = 12,000$ MeV/cm, respectively. That is,

$$k_{\text{hip}} = \begin{cases} 1 & \text{for } dE/dx \leq 969 \text{ MeV/cm,} \\ 0.000754 \cdot dE/dx + 0.269 & \text{for } 969 < dE/dx < 12,000 \text{ MeV/cm,} \\ 9.3172 & \text{for } dE/dx \geq 12,000 \text{ MeV/cm.} \end{cases} \quad (4.29)$$

Originating from the heavy-ion data, the systematic uncertainties associated with the fit are

$$k_{\text{hip}} = \begin{cases} 0.000959 \cdot dE/dx + 0.150 & \text{for } 969 < dE/dx \leq 6,000 \text{ MeV/cm,} \\ 0.000695 \cdot dE/dx + 1.731 & \text{for } 6,000 < dE/dx < 12,000 \text{ MeV/cm,} \end{cases} \quad (4.30)$$

and

$$k_{\text{hip}} = \begin{cases} 0.000724 \cdot dE/dx + 0.101 & \text{for } 969 < dE/dx \leq 6,000 \text{ MeV/cm,} \\ 0.000525 \cdot dE/dx + 1.293 & \text{for } 6,000 < dE/dx < 12,000 \text{ MeV/cm,} \end{cases} \quad (4.31)$$

from its upper and lower variations. The implementation of Equation 4.30 and Equation 4.31 (see Figure 4.8) results in a significant improvement in the simulation, as shown in Figure 4.9. The fact that R_{MC}/R_{Exp} approaches unity as dE/dx becomes larger indicates better modelling of the measured charges or energy due to the effect of charge recombination.

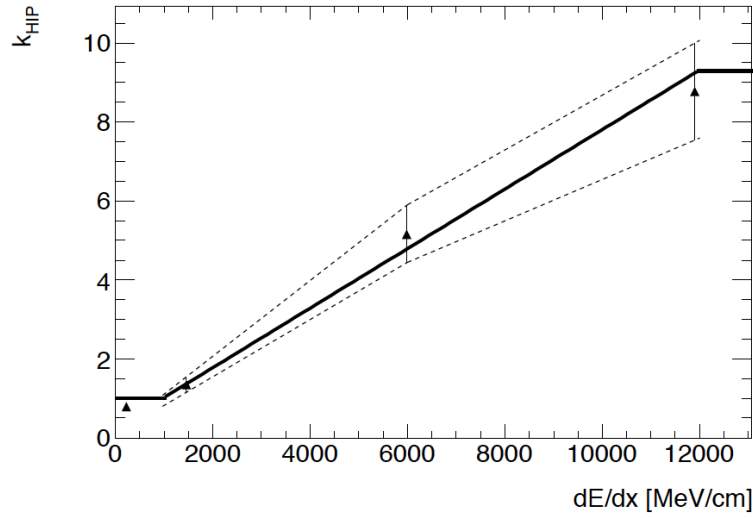


Figure 4.8: HIP correction to Birks' Law (solid line) and its upper and lower variations (dashed lines) as a function of dE/dx for heavy ions [186].

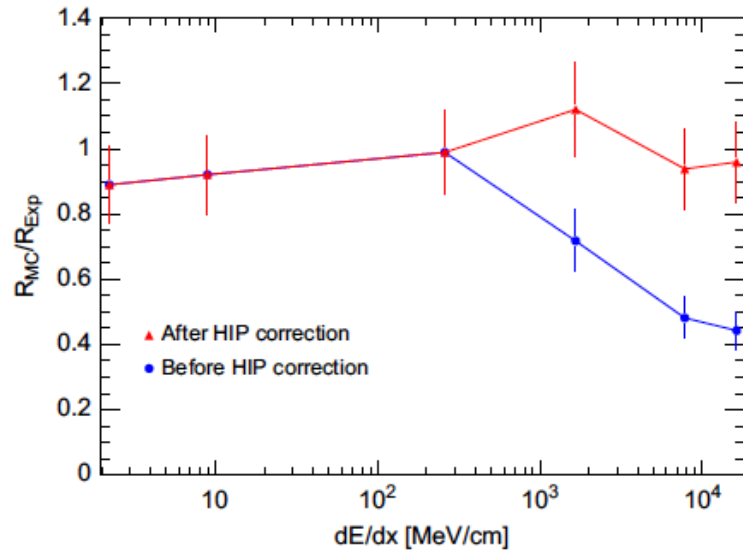


Figure 4.9: Performance (R_{MC}/R_{Exp}) of the ATLAS implementation of Birks' Law as a function of dE/dx for heavy ions [186].

Chapter 5

Highly Ionizing Particle Search Strategy

This chapter is devoted to the methodology of the highly ionizing particle (HIP) search, featuring a Monte-Carlo (MC)-driven characterization of the HIP signal in reconstructing and selecting candidate events. Background rejection in data is balanced against the candidate selection efficiency, via a fraction of the full 2015 – 2018 dataset in conjunction with a set of simulated HIP signal samples. Efforts to that end encompass presorting analysis data via a software framework and fine tuning the event-selection criteria.

5.1 Monte Carlo Samples

This search targets spin-0 and spin- $\frac{1}{2}$ HIP production via the photon-fusion and the Drell-Yan mechanisms. All four production models consider HIPs with masses of 200, 1000, 1500, 2000, 2500, 3500, 4000 GeV, and seven charges, which include $|g| = 1g_D$ and $2g_D$ for monopoles, as well as $|z| = 20, 40, 60, 80, 100$ for HECOs. These masses and charges were chosen to ensure the selection of HIPs at the trigger level, whose efficiency is shown in

Section 5.3.2. The mass range also follows from the collision energy, which diminishes the probability of producing HIPs with mass greater than 4000 GeV.

The ATLAS Monte Carlo (MC) production system manages the production of MC samples on the Worldwide LHC Computing Grid (WLCG), as discussed in Section 3.5. I requested MC samples from this system to investigate detector response and selection efficiencies. Since the HIP interaction with detector material is model-independent, any efficiency variation, for the same charge-mass combination, is caused by kinematic differences between the signal models. Therefore, except for Drell-Yan spin- $\frac{1}{2}$ HIPs, which are fully simulated and reconstructed for the study of detector response, generator-level samples are used for the alternate models, whose efficiencies are sampled from model-independent efficiency maps of fully simulated and reconstructed particle-gun samples [187]. The samples are produced in different subcampaigns that model the LHC operation parameters and ATLAS detector conditions of different data-taking periods, i.e., MC16a models 2015 – 2016, MC16d models 2017, and MC16e models 2018.

5.1.1 Particle Gun

Primarily used in studying the alternative models, the HIP particle-gun configuration framework was originally developed by former members of the analysis to sample η , ϕ , and kinetic energies (E_{kin}) of the HIPs from uniform distributions. I ported the tool to be compatible with the latest release of the ATLAS Athena framework (see Section 3.4). I was also responsible for validating the samples it produced for all charges and masses. Configured by the tool, these particle-gun samples contain HIPs populated equally in the ranges of $-2 < \eta < 2$, $-\pi < \phi < \pi$, and $10 \text{ GeV} < E_{\text{kin}} < 4000 \text{ GeV}$, respectively.

5.1.2 Benchmark Models

After creating new benchmark models (see Chapter 2), I submitted for inclusion in future Athena releases the code of these HIP models, as they were needed for the production. MC samples for each of the four benchmark processes were produced individually with MADGRAPH5_aMC@NLO (MadGraph) [147]. I derived the p_T cuts from the particle-gun trigger efficiencies (see Section 5.3.2) and applied them as cuts on the HIP transverse momenta (p_T) at the generator level to ensure no CPU time is wasted simulating HIPs that are unlikely to penetrate the Electromagnetic Calorimeter and fire the trigger. I maintained the models for proper integration into the rest of the ATLAS production chain as follows. The MadGraph output files for parton-level simulation, i.e., the Feynman diagrams of the Drell-Yan and photon-fusion mechanisms, interface to the PYTHIA8 generator [169] for showering and hadronization. With the A14 tune [188] that describes the underlying events, the Drell-Yan production in the hard collision exploits the NNPDF23LO parton distribution function [150], while the photon-fusion based hard collision is simulated using the LUXqed photon parton distribution function for its better measurement of the photon momentum [151, 152].

5.2 Analysis Framework

The MC simulated events and collision data are processed by an analysis framework that I helped develop. The data formats and the key features of the analysis framework will be described.

5.2.1 Data Flow

The data undergo multiple conversions, as shown in Figure 5.1. The real data takes a simpler route that begins directly from reconstruction to process detector response into

byte-stream RAW data. The data reconstruction is performed at Tier-0 (T0) on a subset of the full RAW dataset that contains events selected by a HIP filter. The RAW data comprises any event that has fired at least one trigger. The HIP filter checks if the HIP trigger bit is true in any given event, and saves the event into the Derived Event Summary Data (DESDM.EXOTHIP) format. On the other hand, the simulated events requires a full chain of processing, from particle-gun- or MadGraph-based event generation that produces truth EVNT files to simulation of particle interactions with the detector, then to digitization algorithms that output Raw Data Objects (RDOs). Taking place at Tier-1 (T1), the MC reconstruction processes every RDO event into the Event Summary Data (ESD) format and saves trigger bits, whether or not the event passes any trigger. The RDO is roughly equivalent to the byte-stream RAW data produced from the High-Level Trigger (HLT).

The ESD and the DESDM.EXOTHIP data preserve the TRT hit information as well as the calorimeter cell information, which are crucial for this analysis. Subsequently, these data files are put through a reduction framework, namely the “ntuple maker code”, which specifies what types of information or collections of data objects should be retained for the analysis. I upgraded this framework and maintained its running in the latest ATLAS Athena environment. This framework consists of an algorithm, which is a custom trimming procedure involving standard ATLAS analysis tools in a class of C++ code. The algorithm runs on each event in the ESD and the DESDM.EXOTHIP datasets under the ATLAS Athena framework via a Python jobOptions file, which steers the execution of the algorithm and configures the implemented standard tools, such as pileup reweighting tools and event quality tools.

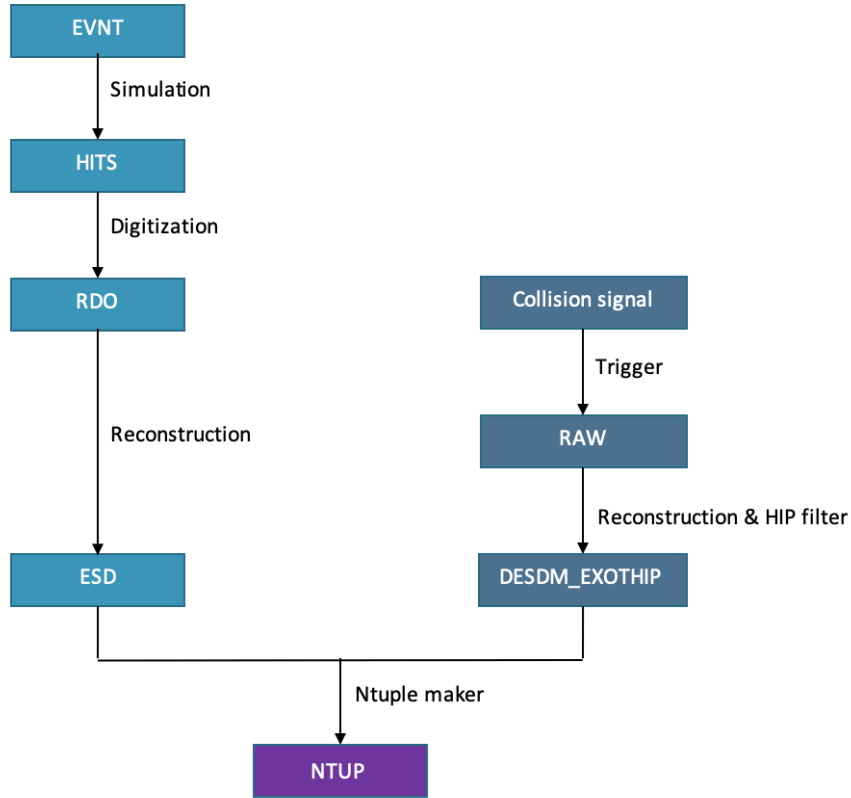


Figure 5.1: HIP data flow: MC (left) and real data (right), both converge to an ntuple (NTUP) format. Note that the MC pathway is depicted for the fully simulated and reconstructed Drell-Yan spin- $\frac{1}{2}$ samples, and thus requires the simulation, digitization and reconstruction steps, which do not apply to the alternate models.

5.2.2 Pileup Reweighting

At the analysis level, the pileup reweighting procedure adjusts the pileup distribution in the MC simulation to match with that in the data, as the MC production parameters are typically frozen prior to data taking.

One difference between these pileup distributions is of vital importance in understanding this procedure. As has been noted in Section 3.4, the MC signal events in a lumiblock are assigned the same amount of pileup regardless of the collision bunch crossing they belong to. In real data, for every measurement of instantaneous luminosity during a lumiblock, the

number of inelastic proton-proton interactions per bunch crossing is distributed according to a Poisson distribution, hence, one takes the mean value of this distribution and records it as the average number of interactions per bunch crossing. This measure of pileup, denoted as $\langle\mu\rangle$, is an average over multiple bunch crossings in a given lumiblock that is longer than the subdetector readout windows. Therefore, $\langle\mu\rangle$ is especially useful in situations where the pileup fluctuation from bunch to bunch is irrelevant, such as the monitoring of data-taking conditions. On the other hand, $\langle\mu\rangle$ proves less appropriate for physics studies, where each data event belongs to a collision bunch whose pileup level is unique. This realization brings about the notion of μ , as a “per bunch crossing” quantity to describe the pileup level, for physics analyses that are more sensitive to pileup fluctuation from bunch to bunch than that from lumiblock to lumiblock. A comparison of μ and $\langle\mu\rangle$ is given in Figure 5.2.

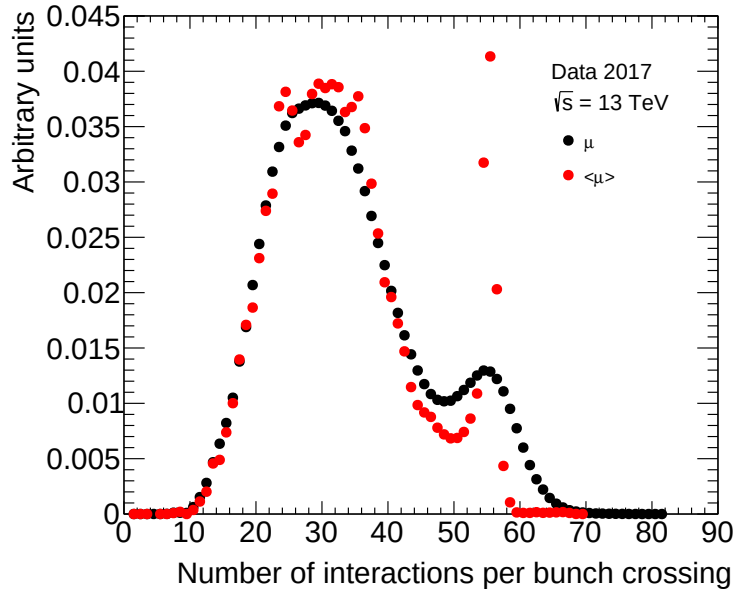


Figure 5.2: Pileup in 2017 data represented by μ and $\langle\mu\rangle$.

In order to reflect the true pileup conditions as recorded in data, I configured the pileup reweighting tool that defines and applies weights to MC samples event-by-event as a correction to the MC pileup distributions. Since MC16d, the use of the μ distribution is recom-

mended over the $\langle\mu\rangle$ distribution to represent the pileup conditions in data. For each MC subcampaign, the pileup weight for an event is calculated based on its μ ($\langle\mu\rangle$) bin i as

$$w_i = \frac{\mathcal{L}_i/\mathcal{L}}{N_i/N}, \quad (5.1)$$

where $\mathcal{L}$ is the total integrated luminosity delivered by the LHC in the corresponding data taking period and N represents the number of MC events. The μ ($\langle\mu\rangle$) distribution of a given sample is scaled by these weights. The μ ($\langle\mu\rangle$) distributions from all three MC subcampaigns are merged and normalized to the distribution of the full 2015 – 2018 dataset, as shown in Figure 5.3.

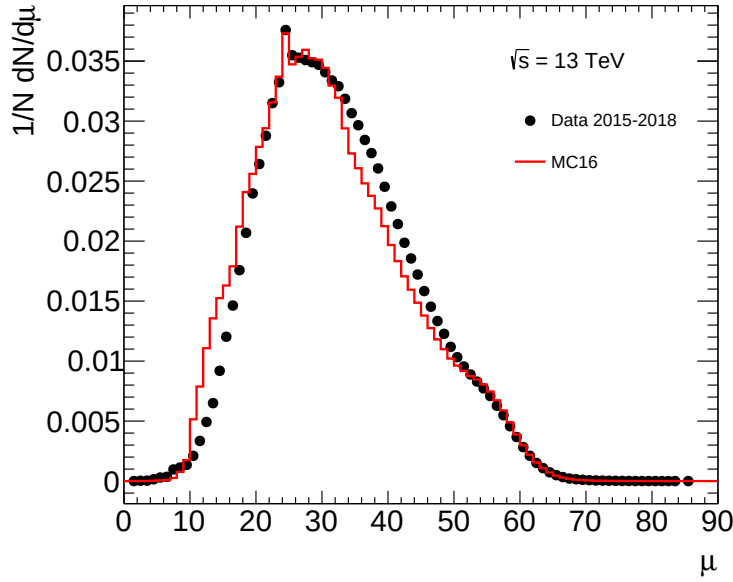


Figure 5.3: Data-MC comparison of pileup.

5.2.3 Event Quality and Basic Cleaning

Although the real collision events must have fired a trigger to be recorded, not all triggered events are recorded with optimal detector conditions. As a result, the collision data undergo

careful scrutiny for the best possible quality of the recorded information with two filters. The first filter rejects data collected with subdetector defects. Compiled centrally, the *Good Runs List* (GRL) records the overall run conditions in all lumiblocks. A dedicated GRL tool is applied in the analysis framework to veto the events that are flagged as suboptimal by the GRL. The second requirement is to ensure the data was neither recorded in time windows with noise bursts in subdetectors nor saved with corrupted event information. Therefore, a set of basic event cleaning procedures is implemented in the analysis framework.

5.3 Event Selection

HIP candidates are selected based on their interaction with the Transition Radiation Tracker (TRT) and the Liquid Argon (LAr) Electromagnetic (EM) Calorimeter described in Chapter 4. In the TRT, the HIPs produce a large number of high-threshold hits, while in the calorimeter, especially in the presampler and the first Electromagnetic Calorimeter layer (EM1), they deposit a large amount of concentrated energies.

5.3.1 Backgrounds

The main sources of background that may disrupt the HIP selection are high-energy electrons and multijets. The high-energy electrons produce TRT high-threshold hits from ionization and from transition radiation; subsequently, they leave compact shower shapes in the EM1 since they do not fully develop electromagnetic cascades until they arrive in the second Electromagnetic Calorimeter layer (EM2). The multijets can fake HIP signals when they overlap in the same TRT straw and leave large and collimated energy depositions in the calorimeter as a result of their high p_T values; they penetrate to the hadronic calorimeter, as observed for some HIPs.

5.3.2 Triggers

The distinctive signature of the HIPs prompted the development a dedicated trigger for the 2012 HIP analysis [67]. The Level-1 (L1) trigger operates with custom electronics in the calorimeters (and the Muon Spectrometer), meaning, a HIP candidate will not be selected unless it arrives at the LAr Electromagnetic Calorimeter and deposits sufficient energy that exceeds a L1 electron-photon (*egamma*) trigger threshold. At the High-Level Trigger (HLT), the custom HIPTRT trigger identifies regions of high ionization densities in the TRT [70, 72, 189, 190].

Trigger Definition

The trigger selection begins with a L1 seed L1EM22VHI, where VHI stands for Variable Hadronic Isolation. The general design of the L1 trigger is discussed in Section 3.2.4. A variable trigger threshold fluctuates around the transverse energy of 22 GeV in response to non-uniform efficiency across the η range of the calorimeter without lowering the overall trigger efficiency. A hadronic veto is applied to calorimeter tower candidates with less than 50 GeV transverse energy when at least 1 GeV is deposited in the hadronic calorimeter, as shown in Figure 5.4, where 1 GeV is an isolation criterion that checks for sufficiently low energy in the isolation ring of the EM clusters. By and large, this design manages to control the trigger rate without losing the trigger efficiency.

Subsequently, the dedicated HIP trigger HLT_g0_hiptrt_L1EM22VHI at HLT selects HIP candidates that have deposited sufficient energy in the Electromagnetic (EM) calorimeter. The trigger algorithm (Trigger/TrigAlgorithms/TrigTRTHighTHitCounter) reconstructs the number $N_{\text{HT, trig}}$ and the fraction $f_{\text{HT, trig}}$ of TRT high-threshold (HT) hits in a wedge around the L1 Region-of-Interest, i.e., an EM calorimeter region with sufficient energy deposit. First, the centre of the wedge is determined through an iterative binning process to select the bin

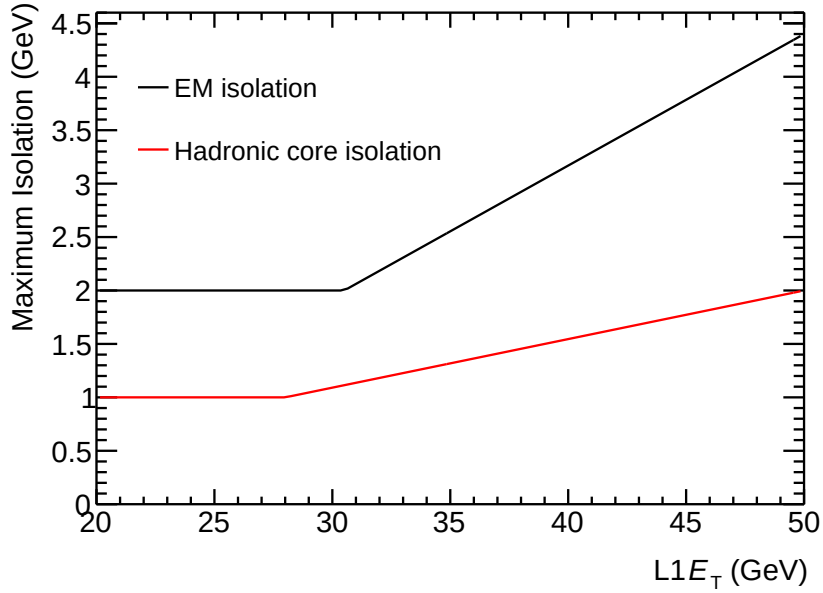


Figure 5.4: L1 calorimeter isolation cuts. The hadronic veto removes calorimeter tower candidates with less than 50 GeV transverse energy when at least 1 GeV is deposited in the hadronic calorimeter (adapted from [191]).

with the highest number of HT hits. The wedge then forms around the centre with a narrower width of 11 mrad in ϕ , which increases from 6 mm to 11 mm as the wedge extends towards the outermost layer of the TRT. The selection criteria $N_{\text{HT, trig}} > 30$ and $f_{\text{HT, trig}} > 0.5$, as well as $|\eta| < 1.7$, control the trigger rate and maintain high signal efficiency. Data selected by the trigger is saved in the DESDM_EXOTHIP format by a dedicated filter.

Signal Trigger Efficiency

The HIP trigger efficiency is dominated by the L1 trigger efficiency, which depends on the HIP transverse momentum p_T , as shown in Figures 5.5 – 5.6 for particle-gun samples. The onsets of the efficiency curves correspond to minimum p_T requirements for the HIPs to pass the trigger threshold, hence, I used these values as low p_T cut-offs in the generator-level simulation of the HIP production models. I derived a p_T cut for each charge-mass combination of a given production model by finding the lower bound of the p_T bin with a

particle-gun trigger efficiency greater than 1%. The trigger efficiency increases with p_T and plateaus at nearly unity, with a few exceptions due to the L1 hadronic veto (see Figure 5.4) for HIPs that deposit between 1 – 2 GeV transverse energies in the hadronic calorimeter but less than 50 GeV in the Electromagnetic Calorimeter. The hadronic veto comes into effect predominantly for $|g| = 2g_D$ monopoles, $|z| = 20, 80, 100$ HECOs, and manifests as deviations from the efficiency plateaus in Figures 5.5 – 5.6, at different p_T regimes.

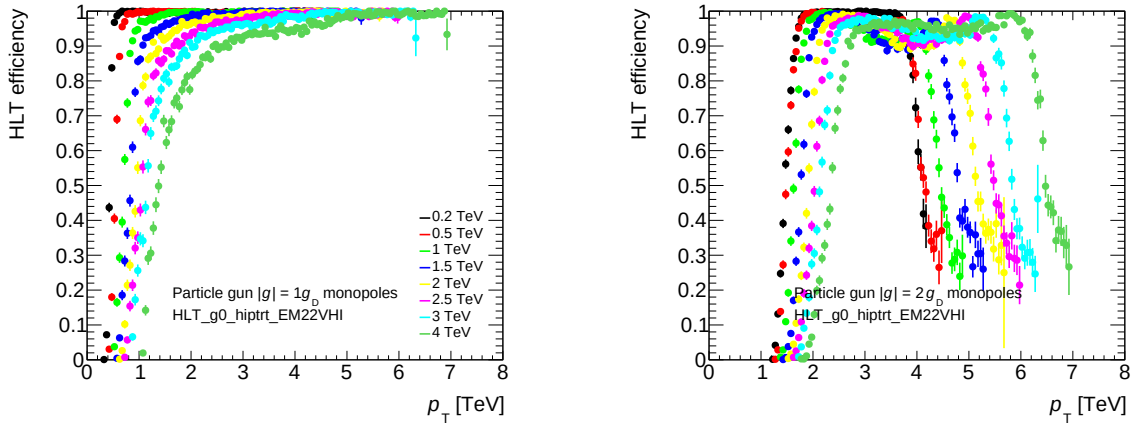


Figure 5.5: HIP HLT efficiencies for monopoles with charges $|g| = 1g_D$ (left) and $|g| = 2g_D$ (right) and various masses, as a function of truth transverse momentum [74]. Particle-gun samples are used for these plots.

The trigger efficiency and its p_T dependence are a manifestation of the HIP dE/dx (Equations 4.11 and 4.12). For the same charge, the efficiency variation is more moderate among different masses for monopoles than for HECOs, as the monopole dE/dx has a weaker dependence on mass than the HECO dE/dx . Nevertheless, HIPs with the same charge show efficiency onsets at higher p_T values if they are heavier.

The HIP efficiency is more sensitive to charge variation than mass, as illustrated in Figures 5.5 – 5.6. First, high-charge HIPs, such as $|g| = 2g_D$ monopoles and $|z| = 100$ HECOs, develop inefficiencies over a wide p_T range; these HIPs are the most ionizing, hence, they require $p_T > 1$ TeV to be triggered and $p_T < 4$ TeV to avoid the L1 hadronic veto. Next, efficiency drops are observed after the onsets and before the plateaus ($2 \text{ GeV} < p_T < 4 \text{ GeV}$)

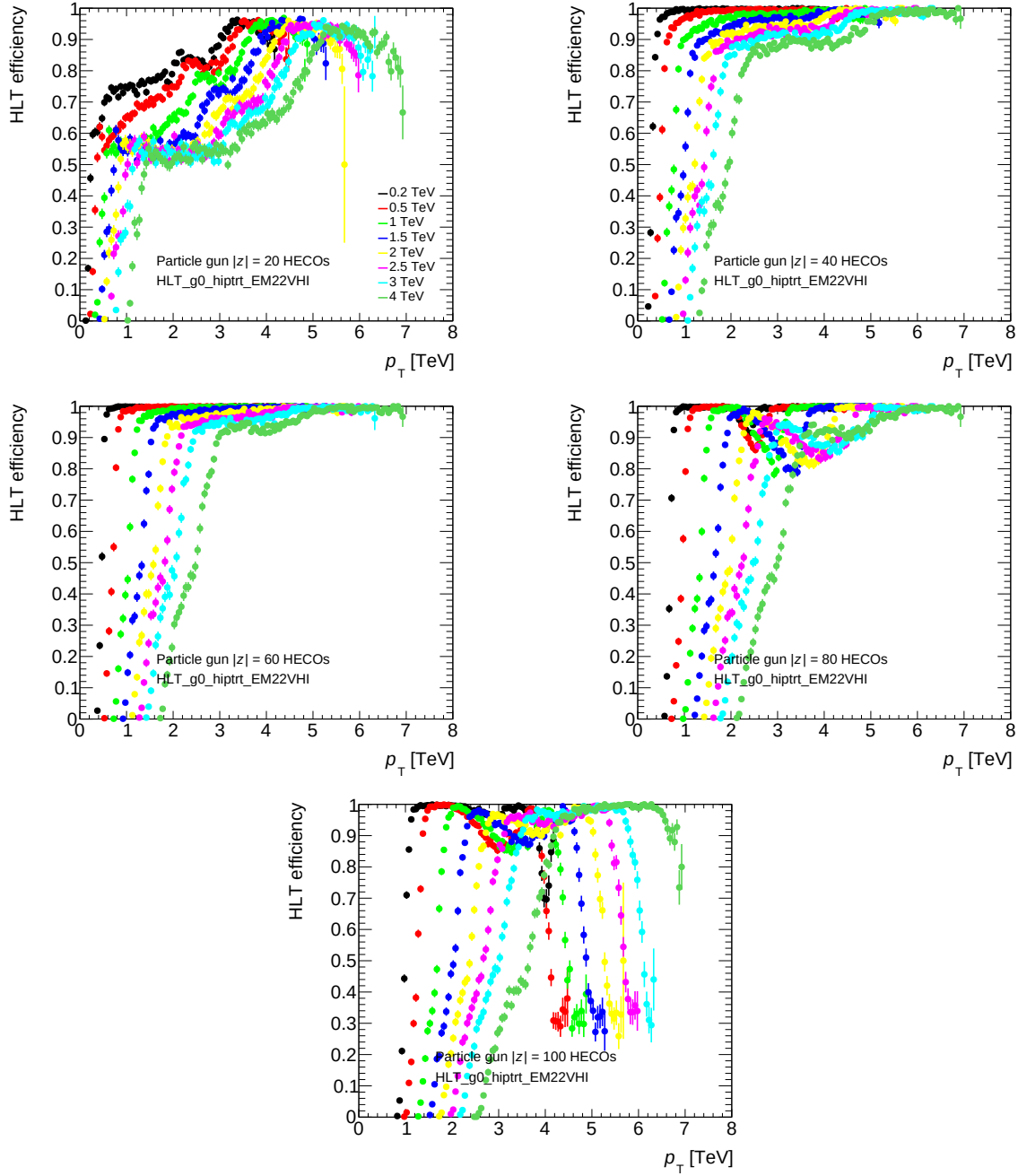


Figure 5.6: HIP HLT efficiencies for HECOs with charges $|z| = 20 - 100$ and various masses, as a function of truth transverse momentum [74]. Particle-gun samples are used for these plots.

for $|z| = 80$ HECOs over a range of masses and $|z| = 100$ HECOs with low masses; this is a sign of energetic HIPs overcoming the L1 hadronic veto effect. Lastly, the $|z| = 20$ HECO

efficiency increases over a wide p_T range after the initial onset but hardly plateaus; these HIPs are the least ionizing, so they remain energetic enough to penetrate into the hadronic calorimeter, which may lead to the firing of the L1 hadronic veto.

The trigger efficiency also reflects the total HIP energy loss. The amount of material before the EM calorimeter is proportional to $1/\sin\theta$ ($1/\cos\theta$) for the barrel (end-caps), meaning the HIP energy loss scales as the truth transverse kinetic energy ($E_{\text{kin}} \sin\theta$) in the barrel. Beyond an optimal range of transverse kinetic energies up to about 4 TeV, the HIP trigger efficiency decreases either because the HIPs fired the hadronic veto or because the HIPs did not deposit enough energy in the calorimeter. Therefore, the HIP trigger efficiency is dependent on the truth pseudorapidity and transverse kinetic energy, which is the founding basis of the HIP efficiency maps to be presented later.

The trigger efficiency is expected to decrease in the presence of pileup, which dilutes the fraction of high-threshold TRT hits with additional low-threshold hits. In Figures 5.7 – 5.8, HIPs with lower dE/dx , i.e., those of higher masses or lower charges, are more susceptible to the pileup effect.

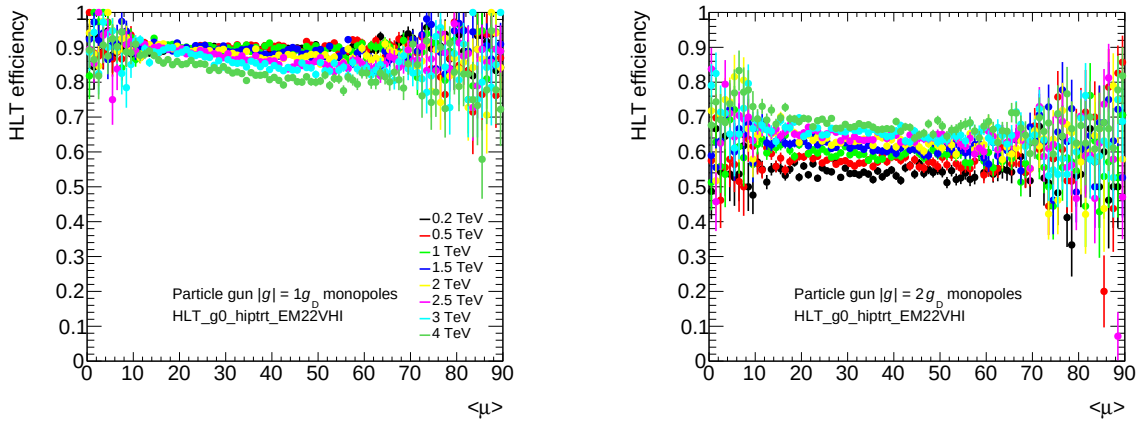


Figure 5.7: HLT efficiencies for monopoles with charges $|g| = 1g_D$ (left) and $|g| = 2g_D$ (right), and various masses, as a function of average number of interactions per bunch crossing ($\langle\mu\rangle$) [74]. Particle-gun samples are used for these plots. These monopoles have kinetic energies less than 4 TeV and p_T ranges given in Figure 5.5.

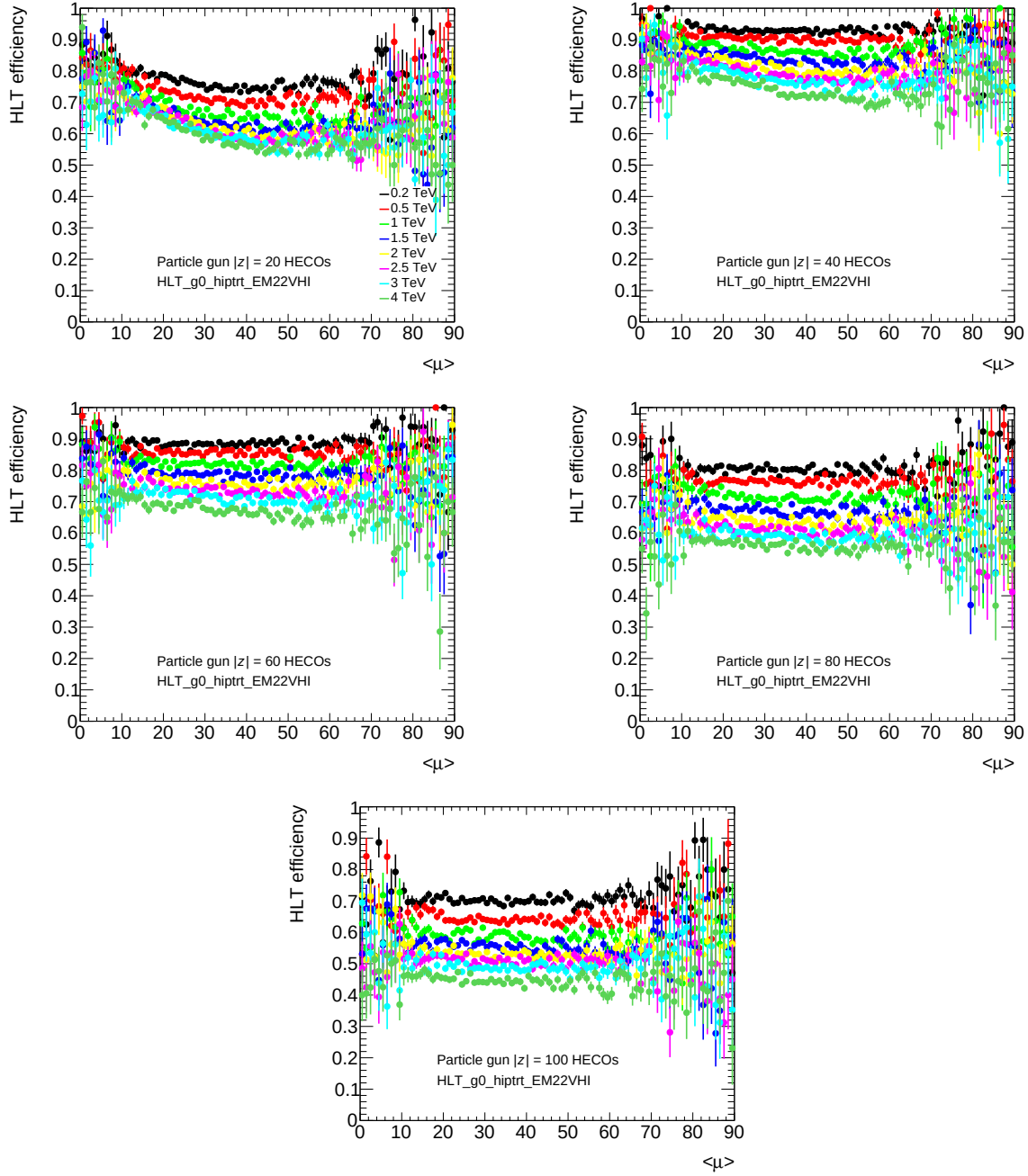


Figure 5.8: HLT efficiencies for HECOs with charges $|z| = 20 - 100$ and various masses, as a function of $\langle\mu\rangle$ [74]. Particle-gun samples are used for these plots. These HECOs have kinetic energies less than 4 TeV and p_T ranges given in Figure 5.6.

Other sources that vary the efficiency will be discussed later as systematic uncertainties.

5.3.3 Preselection

While the trigger is designed to maximize signal efficiency at a reasonable rate, preselection based on calorimeter clusters reduces the size of the data samples for further analysis.

In light of significant upstream energy loss, a HIP candidate in the calorimeter is reconstructed by taking into account the shape and the energy of the calorimeter cluster. Topological clusters, in particular *CaloCalTopoClusters* [192], are used to reconstruct HIP candidates in an energy-density-oriented approach. A *CaloCalTopoCluster* is a group of calorimeter cells that are topologically connected to one another. The clustering algorithm exploits the signal significance of each cell with respect to noise, which is a combination of electronic noise and pileup. The formation of a topocluster is seeded by cells with signal significance above a given threshold. Subsequently, neighbouring cells with signal significance above another (lower) threshold are included in the cluster. Next, all adjacent cells are added to the cluster regardless of their signal significance. The cluster is not allowed to have more than one local energy maximum due to a splitting algorithm. The energy of the resulting *CaloCalTopoCluster* is the sum of all cell energies. The angular coordinates, pseudorapidity and azimuthal angle of the *CaloCalTopoCluster*, is the average of the coordinates of all cells, weighted by their energies.

To slightly tighten the requirement of the calorimeter energy deposition, the clusters are required to have the same minimum transverse energy of 22 GeV as the L1 trigger objects, summing over the presampler and the three layers of the Electromagnetic Calorimeter. These clusters are more compact objects than the L1 trigger objects, hence, this cluster transverse-energy cut refines the selection of calorimeter objects without losing signal efficiency. Additionally, a pseudorapidity cut of $|\eta| \leq 1.375$ is imposed on these clusters to reject the barrel-end-cap transition between $1.375 < |\eta| < 1.52$. The pseudorapidity cut is motivated by the requirement of a well-defined final-selection variable (see Section 5.3.4) and

the necessity to exclude larger $|\eta|$ values that introduce correlation between final-selection variables for background estimation (see Section 5.4).

5.3.4 Final Selection

Once HIP candidates are identified in events that pass the trigger and the preselection, they are further discriminated against background, with the ultimate goal of uncovering a HIP signal for a potential discovery, or in the absence of a signal for limit-setting.

TRT High-Threshold Hits

An algorithm developed in previous iterations of the analysis counts TRT hits due to energy deposition from HIPs and associated δ -rays [70, 72, 189, 190]. For every preselected calorimeter cluster, a wedge of size $\Delta\phi = \pm 0.05$ rad and the highest number of high-threshold (HT) hits is determined around the cluster's azimuthal angle ϕ_{cluster} in the TRT. Centring on this wedge, a narrower wedge of size $\Delta\phi = \pm 0.01$ rad is formed to compute the average ϕ value $\phi_{\text{HT}}^{\text{avg}}$ of all HT hits it contains. This value is used to build a road along the HIP trajectory. In the TRT barrel, where hits are given by their r and ϕ coordinates, the road is rectangular with a width of ± 4 mm that is optimized to contain the hits due to HIP δ -ray production while controlling the amount of pileup hits. In the TRT end-caps, where hits are identified with z and ϕ coordinates, the road is characterized by a fixed angular width, dynamically set by the width of the barrel road. This choice matches the end-cap road to the barrel road at the inner radius (644 mm) of the end-cap. For a ± 4 mm wide barrel road, the end-cap road is ± 0.006 rad (± 4 mm/644 mm) wide in azimuth. The construction of the TRT road is depicted in Figure 5.9.

Within the reconstructed road, the numbers of TRT hits and TRT high-threshold (HT) hits are counted (see Figure 5.10). Albeit present in the background events, the HT hits

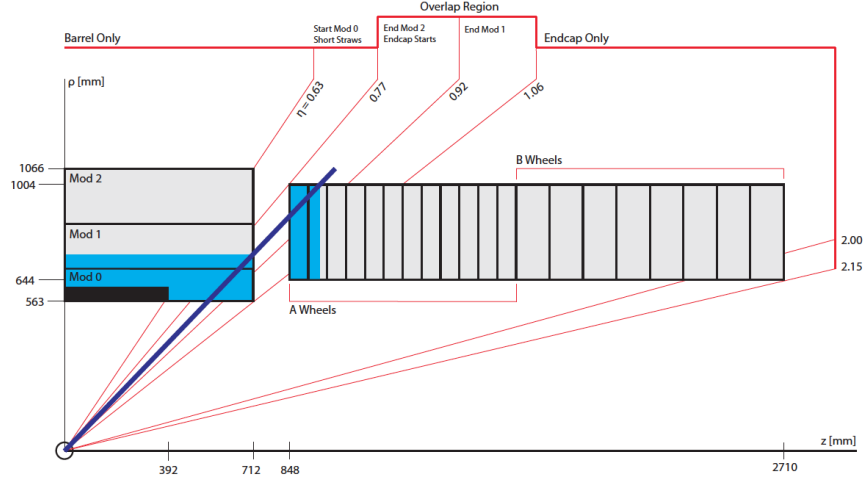


Figure 5.9: Depiction of the TRT road in the r - z plane. The blue line is the HIP trajectory and the blue shaded regions are the straws crossed by the road.

are more prevalent in the signal events from the production of HIPs and associated δ -rays. A minimum requirement of 30 TRT HT hits, expected from the HIP energy deposition described in Section 4.2.3, is followed by the selection of the road with the highest fraction of TRT HT hits (f_{HT} in Figure 5.10) per event. Finally, the selected f_{HT} must be greater than a minimum, which is determined by the technique described shortly.

Lateral Dispersion of Cluster Energy Deposition

The HIP ionization energy loss is mostly contained in the presampler and the first layer of the Electromagnetic Calorimeter (EM1). Hence, the most distinctive signature of a HIP in the EM calorimeter is a cluster with low lateral energy dispersion that is matched to the TRT road with the maximum f_{HT} . The pattern of the HIP energy loss can be described by the lateral arrangement of cells in the cluster. Therefore, a variable is constructed to measure the fraction of energy in the 2, 4 and 5 highest-energy cells in the presampler, EM1, and the second layer of the Electromagnetic Calorimeter (EM2), respectively (see Figure 5.11). The number of cells in each layer is chosen to maximize the signal sensitivity based on

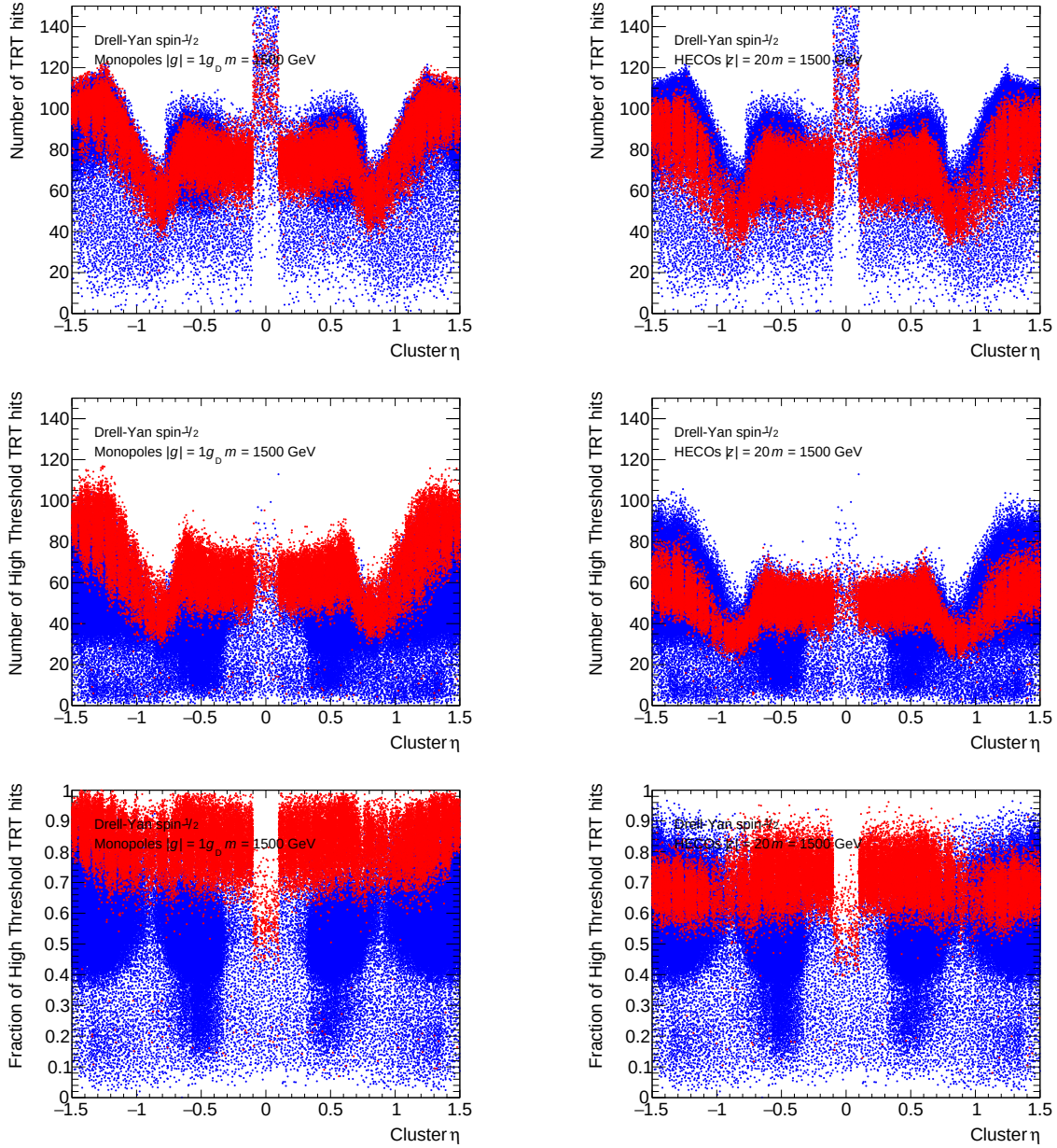


Figure 5.10: The dependence of the number of low-threshold TRT hits, the number of high-threshold TRT hits, and the fraction of high-threshold TRT hits in a ± 4 mm road width as a function of preselected cluster η . The blue points are data. The red points are $|g| = 1g_D$ $m = 1500$ GeV Drell-Yan spin- $\frac{1}{2}$ monopoles (left) and $|z| = 20$ $m = 1500$ GeV Drell-Yan spin- $\frac{1}{2}$ HECOs (right) [74].

previous studies [189]. As shown in Figure 5.12, energy fractions in different calorimeter layers, i.e., w_{pre} , w_{EM1} and w_{EM2} , are separately computed if there is at least an energy of 10 GeV in the presampler, an energy of 10 GeV in EM1, and an energy of 5 GeV in EM2. The variable, w , is computed subsequently as an average over the individual fractions, i.e., $(w_{\text{pre}} + w_{\text{EM1}} + w_{\text{EM2}})/3$ if all three fractions are defined. It is worth noting that the w variable is undefined if the individual fractions are defined for EM2 but not the other two layers. This requirement reduces background contamination of the w variable, as a HIP candidate must penetrate one of these layers to arrive at EM2 and deposit energy. The determination of the cut value is presented next.

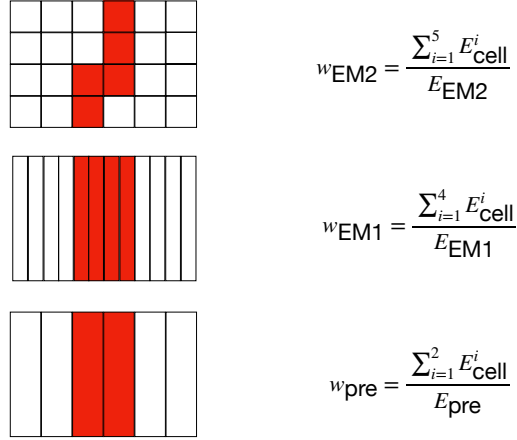


Figure 5.11: The three components of the lateral dispersion of EM cluster energy, where the fractions of energy contained in the highest-energy cells in presampler, EM1 and EM2 are individually calculated.

Cut Optimization

The final-selection cuts strike a balance between signal efficiency and background rejection. This balance can be illustrated with Table 5.1, which tabulates all possible outcomes of the final selection. A perfect cut will count every HIP event as a true signal and every background event as a true background, while a suboptimal cut will lead to a non-zero false-

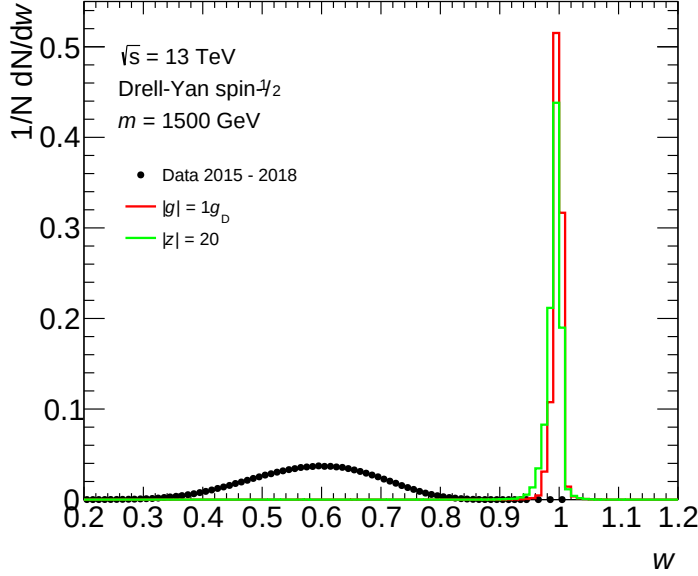


Figure 5.12: Normalized w distributions for data (black), $|g| = 1g_D$ $m = 1500$ GeV Drell-Yan spin- $\frac{1}{2}$ monopoles (red) and $|z| = 20$ $m = 1500$ GeV Drell-Yan spin- $\frac{1}{2}$ HECOs (green).

background rate when a background event fakes a HIP event, or a non-zero false-signal rate when a HIP event is misidentified as a background event.

	Signal output	Background output
Signal input	True signal	False signal
Background input	False background	True background

Table 5.1: Possible outcomes of the final selection, depending on whether or not the cuts reveal the true identity of the input, that is, either signal or background.

Characterizing the outcome distribution, two complementary measures can be exploited to evaluate the performance of a cut. Precision, which is the fraction of selected events that are true signal, namely,

$$\text{Precision} = \frac{\text{True signal}}{\text{True signal} + \text{False signal}}, \quad (5.2)$$

is a measure of the cut's exactness or efficiency at selecting HIP events over fakes. A high

precision rate can be achieved with overly stringent cuts that suppresses the occurrence of false signals but also rejects real HIPs as background. This bias is reflected by Recall, which is the fraction of signal events that are correctly identified, i.e.,

$$\text{Recall} = \frac{\text{True signal}}{\text{True signal} + \text{False background}}. \quad (5.3)$$

A high recall rate translates to a low false-background rate; it measures the completeness of the cut at selecting true HIP events, and thus complements the precision.

In order to determine the f_{HT} cut and the w cut, a balance between the precision and the recall is pursued with the F-measure [193]:

$$F_{\beta} = (1 + \beta^2) \frac{\text{Precision} \cdot \text{Recall}}{\beta^2(\text{Precision} + \text{Recall})}. \quad (5.4)$$

As the harmonic mean of the precision and the recall, the F-measure with $\beta = 1$ is maximized when its two components are equal to each other and decreases otherwise. Different samples yield slightly different cut values, so final cuts are computed as an average weighted by signal efficiencies before the final selection. Different choices of β values in the range of 0 to 2 are also attempted to maximize the F-measure. The resulting cut values fall in the range of $f_{\text{HT}} \in [0.66, 0.88]$ and $w \in [0.85, 0.98]$. The final selection cuts are chosen to be $f_{\text{HT}} \geq 0.77$ and $w \geq 0.93$, as they are around the midrange to provide good statistics for background estimation (to be discussed shortly in Section 5.4) without compromising the background rejection. It should be noted that the data yield after the final selection cuts remains blinded.

5.4 Background Estimation

Although significantly reduced after the event selection, there remain high-energy electrons and high-transverse-momentum multijets that create regions of high ionization den-

sities in the TRT and energy depositions with very low lateral dispersion in the Electromagnetic Calorimeter. These background events are too low in production cross sections to be modelled via Monte Carlo at a reasonable cost of computational resources. Therefore, a data-driven “ABCD” method [194] is used to estimate the residual background, utilizing the two discriminating variables f_{HT} and w . The post-preselection event distribution is presented in a two-dimensional plane, as depicted in Figure 5.13. The cuts on these two variables divide the plane into one signal region A, and three control regions B, C, and D, such that the background distribution in the signal region can be estimated using the control regions.

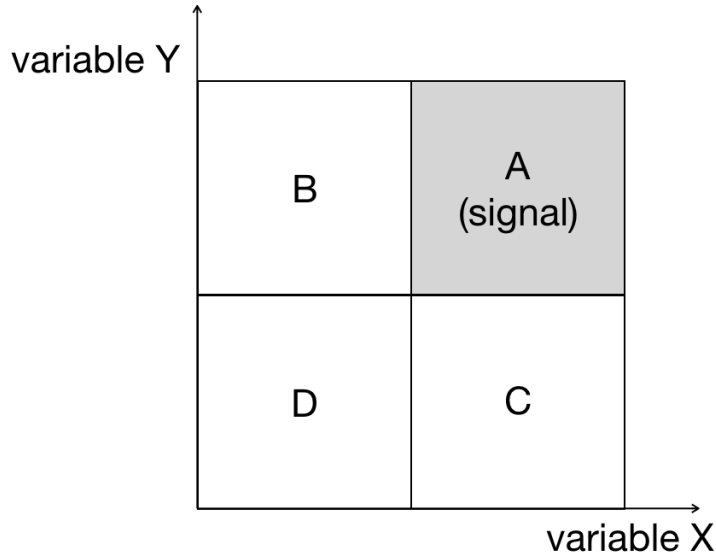


Figure 5.13: Sketch of the ABCD method.

Assuming sufficient data events in the control regions, no signal contamination in the control regions, and no correlation between the two variables, the number of background events in the signal region A can be predicted with

$$\frac{N_C}{N_D} = \frac{N_A}{N_B} \quad \text{or} \quad N_A = N_C \frac{N_B}{N_D}. \quad (5.5)$$

It is obvious that the robustness of this estimate is conditional. First, it will be affected by any correlation between the f_{HT} and w variables; sufficiently low correlation is allowed, however, an uncertainty will need to be assigned to the result. Second, it is important to account for signal leakage, as the background estimate is sensitive to the presence of signals in the control regions as well as in the signal region.

5.4.1 ABCD Method Using Mean Transfer Factor

The correlation between the f_{HT} and w variables is defined by a transfer factor, whose values are computed in the ABCD plane for slices of w as the ratio between the number of events in the region-B portion of a slice and the region-D portion of that same slice, i.e., $\frac{N_B}{N_D}$ in Equation 5.5. If there is no correlation between the two variables, the transfer factor should be uniform such that its numerical values are identical across the w slices regardless of the bin size as long as each bin has sufficient statistics. Although a non-uniform transfer factor is observed in Figure 5.14 as the f_{HT} and w variables are both dependent on η , it should be noted that the correlation between the two variables decreases at high w near the signal region, where an unstable transfer factor biases the background estimation.

This correlation can be regulated as follows. It can be mitigated by excluding η regions with high correlation, i.e., $|\eta| > 1.375$. It can be further reduced by splitting the remaining $|\eta|$ range into subranges, where the transfer factor does not show significant dependence on w . Illustrated in Figure 5.15, the subranges were found to be $0 \leq |\eta| < 0.32$, $0.32 \leq |\eta| < 0.82$, $0.82 \leq |\eta| < 1.375$. For each of these three subranges, I fitted a constant function to the transfer-factor distribution to obtain a mean transfer factor from the fit. The mean transfer factor is quoted with a central value equal to the constant parameter, a statistical uncertainty equal to the error on the constant parameter and a systematic uncertainty equal to the root-mean-square error of the fit.

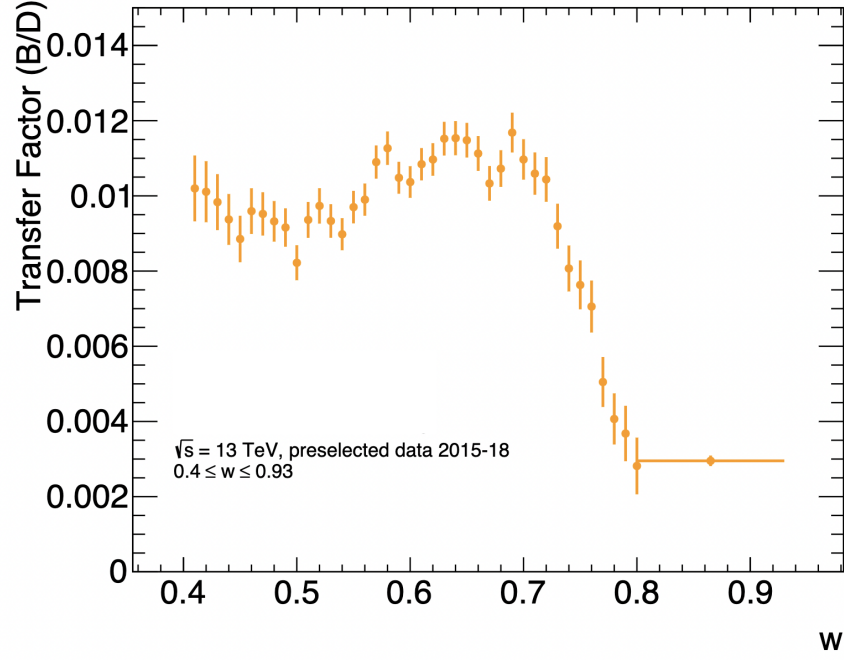


Figure 5.14: Transfer factor as a function of w [74].

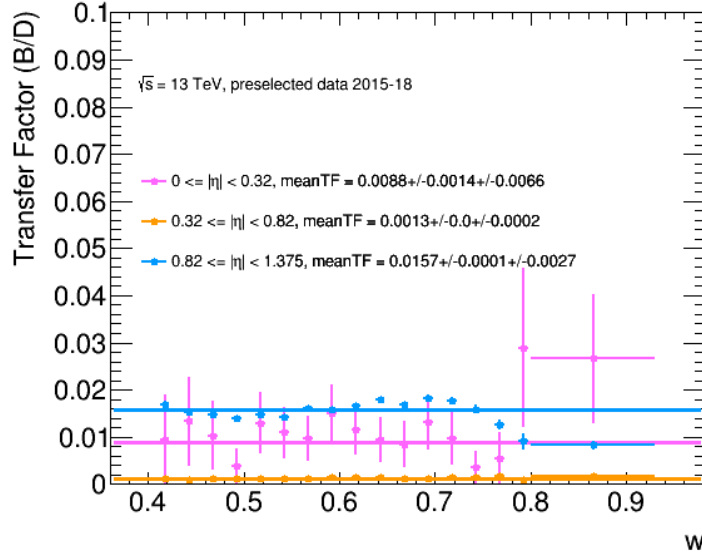


Figure 5.15: Transfer factor as a function of w in three $|\eta|$ slices fitted to zero-degree polynomials. The mean transfer factor is 0.0088 ± 0.0014 (stat) ± 0.0066 (sys) in $0 \leq |\eta| < 0.32$, 0.00129 ± 0.00005 (stat) ± 0.00021 (sys) in $0.32 \leq |\eta| < 0.82$, and 0.0157 ± 0.0002 (stat) ± 0.0027 (sys) in $0.82 \leq |\eta| < 1.375$.

The amount of background in region A can be estimated as the product of the mean transfer factor and the number of background events in C. Since the original ABCD plane is now decomposed into three ABCD planes, one per $|\eta|$ subrange, the background estimate is readily available by computing the background associated with each of these planes and subsequently summing over all three planes.

This method must be validated before it is applied to the data. To this end, subregions are formed in the control regions B and D, as depicted in Figure 5.16. Among the subregions B', B'', D' and D'', B'' is the closest to the signal region, therefore, the background behaviour in B'' can be extrapolated to the signal region A. Two important tests of validation are the variable correlation in Figures 5.14 and 5.15, as well as closure by an equation similar to Equation 5.5:

$$\frac{N_{D''}}{N_{D'}} = \frac{N_{B''}}{N_{B'}}, \quad \text{or} \quad N_{B''} = N_{D''} \frac{N_{B'}}{N_{D'}}, \quad (5.6)$$

where $N_{D''}$, $N_{D'}$, $N_{B''}$, $N_{B'}$ are the numbers of data events in regions D'', D', B'' and B'.

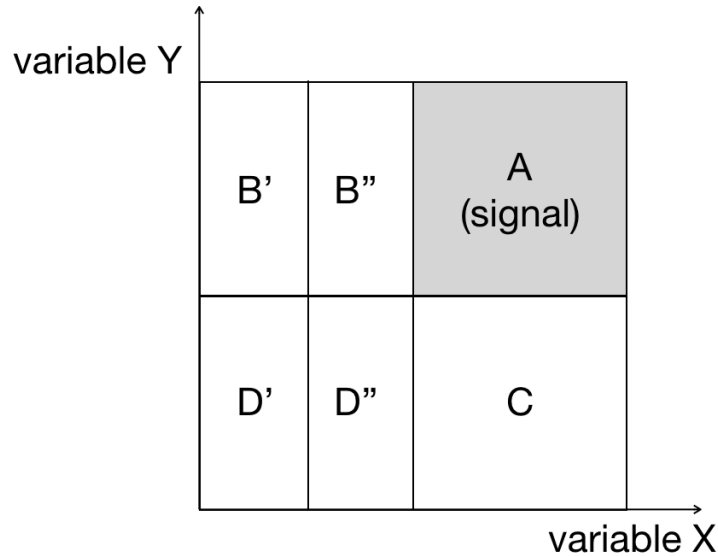


Figure 5.16: Formation of subregions in control regions B and D for the validation procedure.

5.4.2 ABCD Method Using Maximum Likelihood Fit

The method described in Section 5.4.1 is predicated on the assumption that there is no signal leakage into the control regions, i.e., all events observed in the control regions B, C, and D are background events. However, a better assumption can be made with a theoretical prediction of signal leakage whose contribution can be subtracted from the observed events in the control regions, i.e., N_i . Therefore, the true background is

$$N_i^{\text{bkg}} = N_i - N_i^{\text{sig}}, \quad \text{for } i = B, C, D. \quad (5.7)$$

The signal leakage effect that results in non-zero N_i^{sig} can be studied with signal samples. For each sample, the leakage can be quantified as the ratio of signal events in a control region and the signal region:

$$f_i = \frac{N_i^{\text{sig}}}{N_A^{\text{sig}}}, \quad \text{for } i = B, C, D. \quad (5.8)$$

When dealing with significant signal leakage, a suitable remedy is to extract the background estimate from a simultaneous fit to signal and control-region background to better predict background in the signal region. Constructed from the estimated rates of signal processes in all ABCD regions, all sources of systematic uncertainties and the data in the control regions, the probability of finding a certain number of events is modelled by a Poisson distribution for the number of observed candidates and a Gaussian distribution for the various uncertainties that affect this observation. This probability can be interpreted as a likelihood function, whose maximum occurs at a certain number of the background events in the signal region. The HISTFITTER program [195] is exploited to search for the maximum likelihood and the background estimate in the signal region.

As the first step towards describing the probability to observe the real data, the event count can be modelled separately for the $|\eta|$ -sliced signal region and control regions on the

ABCD plane, to account for the signal leakage of the fully simulated spin- $\frac{1}{2}$ Drell-Yan HIPs. Every $|\eta|$ slice enters this model as a channel, whose event count depends on the background and the predicted signal rates with associated uncertainties. For the spin-0 Drell-Yan HIPs and the photon-fusion pair-produced HIPs, it is not possible to either estimate the signal leakage given the statistics of the efficiency maps in the control regions (discussed later in Section 6.3.3) or slice the signal region as the cluster- $|\eta|$ information is not available for these generator-level HIPs. Therefore, the event count constructed from the background and the signal rates of these alternative production models is only modelled for one channel, i.e., the signal region. Although the MC predicts the signal event yields, the expected number of signal events in data is unknown and thus is allowed to float in the model through the parameter of signal strength, μ , that is defined as the predicted signal rate normalized to the actual rate. For each channel i , the expected signal yield s_i and the expected background b_i are separately modelled. Except for channels in the signal region A, where the expected signal $s_i = \mu$, the other channels are subject to non-trivial signal leakage l_i , thus the expected signal yield is modelled as $s_i = \mu l_i$ while $l_i = 1$ for the signal-region channels. The number of background in each signal-region channel is unknown. However, it can be estimated by the transfer factor t and the known background in the control regions, since it follows from Equation 5.5 that $b_A/b_C = t \cdot b_B/b_D$ for each $|\eta|$ slice, where t has been determined in Section 5.4.1. Putting these together, the expected number of events for each channel i is parametrized as $\nu_i = \mu s_i + b_i$. A Poisson distribution that describes the probability of finding n_i data events while ν_i is expected reads

$$\text{Pois}(n_i|\nu_i) = \frac{e^{-\nu_i} \nu_i^{n_i}}{n_i!}. \quad (5.9)$$

This probability distribution is a likelihood function that provides a set of values for the counting model's parameters, i.e., μ , b_i as in ν_i , consistent with the observed data. In other

words, with n^i data events in the i -th bin, the likelihood function

$$\mathcal{L}(\mu, b_i) = \text{Pois}(n_i | \mu s_i + b_i) \quad (5.10)$$

is capable of providing estimates of the model parameters with a range of uncertainty that is constrained by the observed data, as for a given parameter there exists a value that maximizes the likelihood function. This value, known as the maximum likelihood estimate of the parameter, can be determined through the negative logarithm of the likelihood function (log-likelihood) in measures of model fit. The negative log-likelihood is computed iteratively for an initial guess of the model parameters and for successive updates to that guess, with a lower value implying a better fitting model. This process converges if within finite iterations the difference between successive values of the parameters becomes sufficiently small such that the minimum has been achieved. The signal strength and expected background in all bins are obtained from the maximum likelihood fit. The MINUIT program [196] is exploited in the parameter estimation with its built-in minimization algorithms.

It should be noted that Equation 5.10 is a simplified model. Systematic uncertainties (see Section 6.2) and statistical uncertainties (see Section 6.1) affect the expected signal yield, and thus should be accounted for by additional model parameters. These parameters are subject to constraints that allow them to float according to certain probability distributions. In particular, the modified statistical model yields a smaller likelihood if a parameter finds a value far from its expectation. Thus, the likelihood function takes the following form on account of the uncertainty parameters

$$\mathcal{L}(\mu, \boldsymbol{\theta}) = \prod_i^{\text{bins}} \text{Pois}(n_i | \mu s_i(\boldsymbol{\theta}) + b_i(\boldsymbol{\theta})) \times G(\boldsymbol{\theta}), \quad (5.11)$$

where $G(\boldsymbol{\theta})$ is a product of uncertainty distributions that follow Gaussian statistics. The

width of each uncertainty distribution is set by the magnitude of the uncertainty itself. As a consequence, higher uncertainties suppress $G(\boldsymbol{\theta})$ and thus tend to suppress the likelihood function in Equation 5.11, even though they may increase the Poisson probability $\text{Pois}(n_i|\mu s_i(\boldsymbol{\theta}) + b_i(\boldsymbol{\theta}))$.

5.5 Hypothesis Testing

The ATLAS Collaboration requires that all details discussed above in this chapter be finalized before the analysis team can request to unblind the signal region. The unblinding allows the team to determine if there is a statistically significant amount of events in excess of the theoretical prediction. If no discrepancy is observed when comparing the number of candidates in the signal region to the Standard-Model prediction, one sets exclusion limits to the signal models considered.

The first step of the statistical analysis is to test the observed data along with estimated background, against two hypotheses. Known as the background (b)-only hypothesis, the null hypothesis states that there is no sign of new physics, and thus the observed data are fully attributed to Standard-Model processes. Alternatively, the signal + background ($s+b$) hypothesis states that the observed events include not only Standard-Model processes but new physics processes, i.e., the production of HIPs via Drell-Yan and photon-fusion mechanisms. Each hypothesis is rejected based on any inconsistency between the observation and the estimation. The level of inconsistency is reflected by the signal strength μ (explained in Section 5.4.2), with $\mu = 0$ corresponding to the background-only hypothesis and $\mu = 1$ the signal + background hypothesis.

In practice, the compatibility between a hypothesis and the data is measured by a test statistic, q_μ , for a given μ . The test statistic is constructed as the logarithm of the ratio

between two likelihoods [197] as

$$q_\mu = -2 \ln \left(\frac{\mathcal{L}(\mu, \hat{\boldsymbol{\theta}}_\mu)}{\mathcal{L}(\hat{\mu}, \hat{\boldsymbol{\theta}})} \right), \quad (5.12)$$

where $\mathcal{L}(\hat{\mu}, \hat{\boldsymbol{\theta}})$ is the likelihood with best fit values $\hat{\mu}$ and $\hat{\boldsymbol{\theta}}$, while $\mathcal{L}(\mu, \hat{\boldsymbol{\theta}}_\mu)$ is the likelihood maximized for a given μ . The ratio of the two likelihoods is a positive quantity between 0 and 1, thus it follows that the test statistic $q_\mu \geq 0$. The observed value q_μ^{obs} can be computed for data. Large q_μ^{obs} is expected when μ is far from the value specified by the hypothesis and gives evidence against the hypothesis.

The significance of disagreement between the data with a particular hypothesis is computed as the p -value (p_μ)

$$p_\mu = \int_{q_\mu^{\text{obs}}}^{\infty} f(q_\mu|\mu) dq_\mu, \quad (5.13)$$

where q_μ^{obs} is the value of the test statistic q_μ observed in data and $f(q_\mu|\mu, \theta)$ is the probability distribution function of q_μ^{obs} for a given signal strength (μ) hypothesis. In the case where an excess in data is observed, the p -value, namely, the probability that the background can produced a fluctuation greater than or equal to the excess observed in the data is defined using the test statistic q_μ with $\mu = 0$ as

$$p_0 = \int_{q_0^{\text{obs}}}^{\infty} f(q_0|\mu = 0) dq_0. \quad (5.14)$$

To claim a discovery, small p -values are required such that the observed data gives strong evidence against the null hypothesis.

If no significant excess is observed, exclusion limits on the model parameters can be estimated within a certain degree of probability using the modified frequentist method called

“CL_s” [198, 199]. The confidence level CL_s is computed as

$$CL_s(\mu) = \frac{p_\mu}{1 - p_0}, \quad (5.15)$$

where p_μ for $\mu = 1$ and p_0 are defined in Equations 5.13 and 5.14, respectively. The hypothesis of μ is excluded at a confidence level of 95% if $CL_s < 0.05$.

Chapter 6

Highly Ionizing Particle Search Result

This chapter presents the result of the 2015 – 2018 HIP search, performed in accordance with the analysis strategy presented in the previous chapter. Applying the event selection, the HIP signal efficiencies are determined along with the uncertainties of associated systematic effects. Subsequently, the modeling of background and signal event counts comes under close scrutiny. After carefully estimating the expected signal event counts, the signal region is unblinded to reveal the true number of signal events. These event counts are interpreted as limits on the benchmark production models.

6.1 Signal Efficiencies

The signal efficiency ϵ , evaluated entirely from simulation, is defined as the ratio between the number of reconstructed Monte-Carlo (MC) events and the number of MC-generated events:

$$\epsilon = \frac{N_{MC}^{\text{rec}}}{N_{MC}^{\text{tot}}}. \quad (6.1)$$

The efficiency ϵ may be factorized as a product of three efficiency terms:

$$\epsilon = \epsilon_{\text{trig}} \times \epsilon_{\text{reco}} \times \epsilon_{p_T}. \quad (6.2)$$

The trigger efficiency ϵ_{trig} is calculated as the number of triggered candidate events divided by the number of simulated events. The efficiency to reconstruct HIP candidates via selection criteria set out earlier, ϵ_{reco} , is given by the number of selected events divided by the number of triggered events, where an event is accepted as long as one HIP candidate passes all selection criteria. The p_T -cut efficiency, ϵ_{p_T} , quantifies the loss of phase space in the presence of such a cut (see derivation in Section 5.3.2) and is the theoretical cross section with the p_T cut divided by the theoretical cross section without the p_T cut.

On the grounds that the simulation well models the data, the transformation of the number of reconstructed simulated candidates into a signal selection efficiency ϵ allows one to estimate the selection efficiency in the data,

$$\epsilon = \frac{N_{\text{data}}^{\text{rec}}}{L \cdot \sigma}, \quad (6.3)$$

where $N_{\text{data}}^{\text{rec}}$ is the number of reconstructed HIPs (excluding the estimated background events) in the data, L is the integrated luminosity of the data.

The event selection criteria are highly efficient for signal events that pass the HIP trigger. The most significant loss in efficiency, i.e., 50% on average, is incurred at the trigger level, in particular from the First-Level (L1) trigger component, for hadronically-vetoed low-charge HIPs and for high-charge HIPs that stopped before the EM Calorimeter (see Section 5.3.2). The event selection efficiencies for Drell-Yan spin-0 HIPs are the highest due to their kinematic properties (see Chapter 2).

Since all models except for Drell-Yan spin- $\frac{1}{2}$ are not fully simulated, the computation of

event selection efficiencies will be discussed separately for the Drell-Yan spin- $\frac{1}{2}$ model and the alternative models. All event selection efficiencies, assumed to be binominally distributed, are quoted with statistical uncertainties due to the limited size of the simulated samples.

6.1.1 Fully Simulated Efficiency

The signal selection efficiencies for all spin- $\frac{1}{2}$ Drell-Yan HIPs are tabulated in Appendix A Table A.1. The highest signal efficiency of 25% is achieved for Drell-Yan spin- $\frac{1}{2}$ monopoles with mass 1500 GeV and charge $|g| = 1g_D$.

6.1.2 Extrapolated Efficiency

The alternative model efficiencies require event-selection efficiency maps of fully simulated and reconstructed particle-gun HIPs in the truth transverse kinetic energy, E_T^{kin} , versus the absolute value of truth pseudorapidity, $|\eta|$, space. While the choice of $|\eta|$ is dictated by the detector layout, the choice of E_T^{kin} is a direct consequence of the preselection of EM topoclusters of $|\eta| \leq 1.375$. The amount of material before the EM calorimeter barrel region ($|\eta| < 1.375$) is approximately inversely proportional to $\sin \theta$ while before the end-cap region ($|\eta| > 1.52$) it is approximately inversely proportional to $\cos \theta$. Thus the amount of kinetic energy a HIP loses is dependent on its truth $E^{\text{kin}} \sin \theta$ if it arrives at the EM calorimeter and passes the preselection. In order to account for the material in front of the EM calorimeter, efficiency maps are binned in the HIP's truth $E_T^{\text{kin}} = E^{\text{kin}} \sin \theta$ as well as truth $|\eta|$. The bin sizes are chosen to have relatively smooth variations of efficiency with truth E_T^{kin} and with truth $|\eta|$. Figure 6.1 is the efficiency map for particle-gun monopoles of mass 1500 GeV and charge $|g| = 1g_D$.

The generator-level HIPs are subsequently exploited to calculate the event selection efficiency for a given model. An efficiency is extracted from the efficiency maps for each

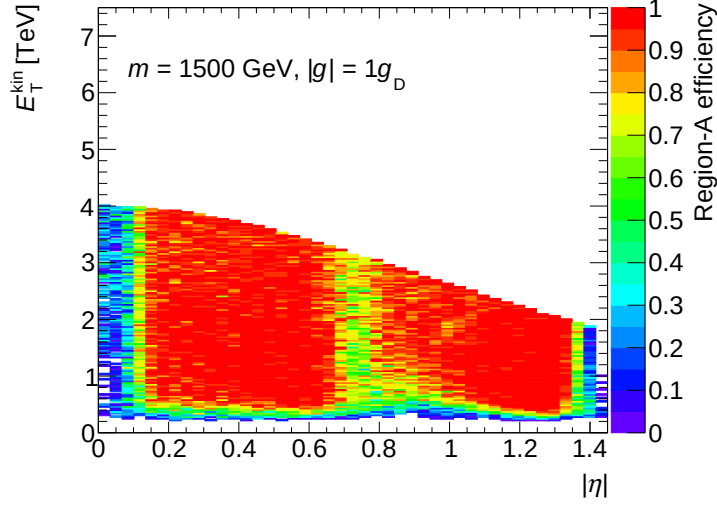


Figure 6.1: Efficiency map of particle-gun monopoles with charge $|g| = 1g_D$ and mass $m = 1500$ GeV.

generator-level HIP given its mass, charge, transverse kinetic energy and absolute value of pseudorapidity. This HIP is passed with a probability equal to the extracted efficiency. This probability is compared against a random number to determine if that event passes the selection. For pair-production models, an event is deemed to pass when at least one of the generator-level HIPs is successfully selected. The extrapolation algorithm loops over all generator-level events of a given model 50 times to attain a binomial distribution of sampled efficiencies. The mean of the distribution is taken as the event selection efficiency, as presented in Appendix A Table A.1 for the Drell-Yan spin-0, photon fusion spin-0 and photon fusion spin- $\frac{1}{2}$ models. The highest signal efficiency of 54% is achieved for Drell-Yan spin-0 monopoles with mass 1500 GeV and charge $|g| = 1g_D$.

6.2 Systematic Uncertainties

Enumerated below are the sources of systematic uncertainties that affect the HIP signal efficiency estimation and are thus included as input to the final results. For instance,

the extrapolation method, as an alternative to full detector simulation, yields slightly different efficiencies. The other sources probe the accuracy with which the pileup simulation and the detector simulation of the TRT and the Liquid Argon (LAr) Electromagnetic (EM) Calorimeter model the HIP efficiency; they are evaluated by varying the parameters germane to the systematic effects in the simulation or at the analysis level. I studied these sources of uncertainties extensively, sometimes in consultation with experts outside the HIP analysis team. I finalized these uncertainties with a systematic component and a statistical component. Each systematic uncertainty is a relative variation of the signal efficiency. The statistical uncertainty in a single variation of the signal efficiency is calculated assuming a binomially distributed efficiency. Unless otherwise specified, the uncertainties tabulated in this section only apply to Drell-Yan spin- $\frac{1}{2}$ HIPs. The full set of uncertainties including those for the alternative models are compiled in [Appendix B](#).

6.2.1 Detector Material

The energy loss and ultimately the signal efficiency of a particle depends significantly on the amount of detector material it traverses, given by the detector geometry, which is a software description of the geometrical layout and material composition for the simulated detector. For highly ionizing particles with large energy losses, the signal efficiency is subject to uncertainty incurred by distorted geometry in the simulation. Therefore, a systematic uncertainty due to material mismodelling is taken into account by considering a detector geometry with more material in front of the Electromagnetic Calorimeter with respect to the nominal geometry. The alternative geometry includes a 10% increase of Pixel (IBL included) and SCT services, a material increase of 7.5% in the Inner Detector and a 5% material increase in the calorimeter presampler. The relative systematic uncertainty is $(9 \pm 4)\%$ on average.

6.2.2 δ -Ray Production Model

Also associated with the HIP dE/dx [40, 41] (see Section 4.1.2), the production of δ -rays affects the energy deposition and reconstruction of the HIP trajectory, thereby affecting the HIP efficiency. First, the dE/dx , given by Equations 4.11 and 4.12, determines the energy deposited in the Liquid Argon (LAr) Electromagnetic (EM) Calorimeter and whether the HIP penetrates into the hadronic calorimeter, dictating the L1 trigger efficiency. Secondly, the dE/dx correlates to the amount of δ -rays produced in the tracker, which create high-threshold (HT) hits in the TRT used in the HLT triggering and offline selection. Lastly, the dE/dx determines the lateral dispersion of the signals in the LAr EM Calorimeter, as parametrized by the w variable used in the offline selection.

The introduction of the systematic effect involves modifying the δ -ray production in the `G4mplAtlasIonisation` class of the custom GEANT4 monopole package (See Section 4.2.2). I introduced a 3% suppression of the production, motivated by the monopole dE/dx uncertainty (see Table 4.1). After a private build for this change in the Athena software framework, I requested particle-gun signal MC samples to compute the alternative efficiencies and the systematic uncertainties.

The HECO ionization is simulated by a different GEANT4 class, which is invoked for ionization by default and thus modification by ATLAS users is not allowed. Therefore, the HECO systematic uncertainties are estimated from the monopole uncertainties, such that HECOs of charges $|z| = 20, 40$ and 60 have the same uncertainties as monopoles of charges $|g| = 1g_D$, while HECOs with charges $|z| = 80$ and 100 are assigned the same uncertainties as monopoles of charges $|g| = 2g_D$. The relative systematic uncertainty is $(2 \pm 0)\%$ on average.

6.2.3 TRT Range Cuts

The HIP ionization and the associated production of δ -rays have large cross sections at low energies. In order to reduce computing resources, the simulation of these processes involves energy cuts that are imposed to forbid the passage of low-energy δ -rays. Although the δ -rays are not explicitly simulated if their kinematic energies are below a threshold, their energy deposition to the detector material remains valid through additional HIP energy losses as the HIPs are propagated. This energy threshold is set by a production cut in range [170] that is specified in the GEANT4 simulation. Generally speaking, the best simulation precision is achieved when the range cuts are kept to their lowest possible values.

The TRT electron range cuts alter the fraction of energy explicitly deposited by the δ -rays relative to that deposited by the HIP and thus ultimately affect the number of TRT hits around the HIP trajectory. As a consequence, these cuts are germane to the HIP efficiency primarily through their effect on the fraction of TRT high-threshold (HT) hits. Different TRT components have different range cuts. For the straw gas mixtures, such as the xenon mixture, the cut is at its optimal value of 30 mm as it suppresses double counting of δ -rays produced in the simulation step by the Photo-Absorption Ionization (PAI) model [179] in the digitization step (see Section 4.2.3). Although the 30 mm cut restricts the δ -rays' kinetic energies to be on the order of 100 keV and higher, it should be noted that variations of this cut do not affect the number of hits produced. The reason lies in the binary readout scheme of the straws (see Figure 3.6), where the same digital signal is recorded for energy depositions above a threshold. The HIPs are so ionizing that their energy depositions are significantly higher than this threshold. That is, if δ -rays are produced in the same straw where a HIP deposits energy, then regardless of their energies they do not change the digital readout. The other range cut is set for the radiator material surrounding the straws to enhance transition radiation (TR). This material contributes to the HIP signal, since it can

be ionized by the HIPs, in which case ionization electrons and δ -rays will travel to adjacent straws, making additional hits. The digital signals of the nearby straws can be biased if the energy deposition by the δ -rays are suppressed due to a large range cut applied to the TR enhancing material. This range cut is set to 1 mm and is sufficiently low, as it ensures that δ -rays are explicitly simulated even if their kinetic energies are significantly below the typical kinetic energy of 1 MeV. Therefore, neither range cut is considered as a source of uncertainty.

6.2.4 TRT Gas Geometry

The modelling of straw gas in simulation plays an important role in the accuracy of the HIP efficiency, as the TRT high-threshold probability depends on if the gas is argon or xenon. In order to achieve a reliable efficiency, the simulation needs to use a gas geometry, i.e., a software description of the straw gas, that is as close as possible to the actual gas used during data taking.

The TRT gas geometry was modified twice during Run 2 in order to replace xenon with argon in leaking straws [200], resulting in three different baselines in data, as shown in Figure 6.2. The first baseline was set in 2015 when argon gas was used in the innermost barrel layer along with end-cap wheel 5 on side A and wheel 3 on side C. The second baseline emerged in 2016, characterized by the use of argon in two barrel layers (0 and 1), two wheels (3 and 5) in end-cap A and three wheels (3, 5 and 8) in end-cap C. The last baseline was benchmarked during 2017 and 2018 due to the change from an argon-based wheel 3 to an argon-based wheel 2 in end-cap C. On the other hand, the simulation implements the same gas geometry as in the 2016 baseline for all years, meaning, the gas was mismodelled by a change of gas between two wheels in 2017 – 2018 and by an excess of argon straws in 2015. Since the change in the 2017 baseline leaves no significant effect on the HIP efficiency, and

the mismodelling of the year 2015 impacted less than 3% of the full dataset, this systematic uncertainty due to the gas mismodeling is deemed negligible.

6.2.5 TRT Occupancy

As far as the TRT is concerned, its last explicit contribution to the HIP efficiency is its occupancy, namely, the average fraction of straws recording low-threshold hits. Any mismodelling of the TRT occupancy in MC may bias the HIP efficiency, because the fraction of high-threshold hits is inversely correlated to the occupancy.

This mismodelling is reflected in all MC events when compared to the corresponding real data, therefore its effect is studied using the decays of Z^0 bosons to electrons and positrons. In order to make the comparison, I requested special data and MC samples to be produced centrally to save the TRT occupancy information. As shown in Figure 6.3, the TRT occupancy is slightly underestimated in the simulation across the entire $\langle\mu\rangle$ range. To assign a systematic uncertainty to the HIP efficiency, the number of low-threshold hits in each HIP signal event is scaled by a $\langle\mu\rangle$ -dependent ratio of occupancy between data and simulation, shown in the lower pad of Figure 6.3. The resulting uncertainty is $(2 \pm 0)\%$ on average.

6.2.6 Calorimeter Arrival Time

The LAr EM Calorimeter is also of primary importance for the HIP measurement. It has a long readout time window that exceeds the 25-ns time gap between two adjacent collision bunch crossings. It identifies particle energy deposition associated with the current collision bunch crossing and compares this energy against the trigger threshold. For any particle to be triggered in the correct collision bunch crossing, it must arrive at the calorimeter before the next bunch crossing. Otherwise, it can produce a signal pulse in the next collision bunch

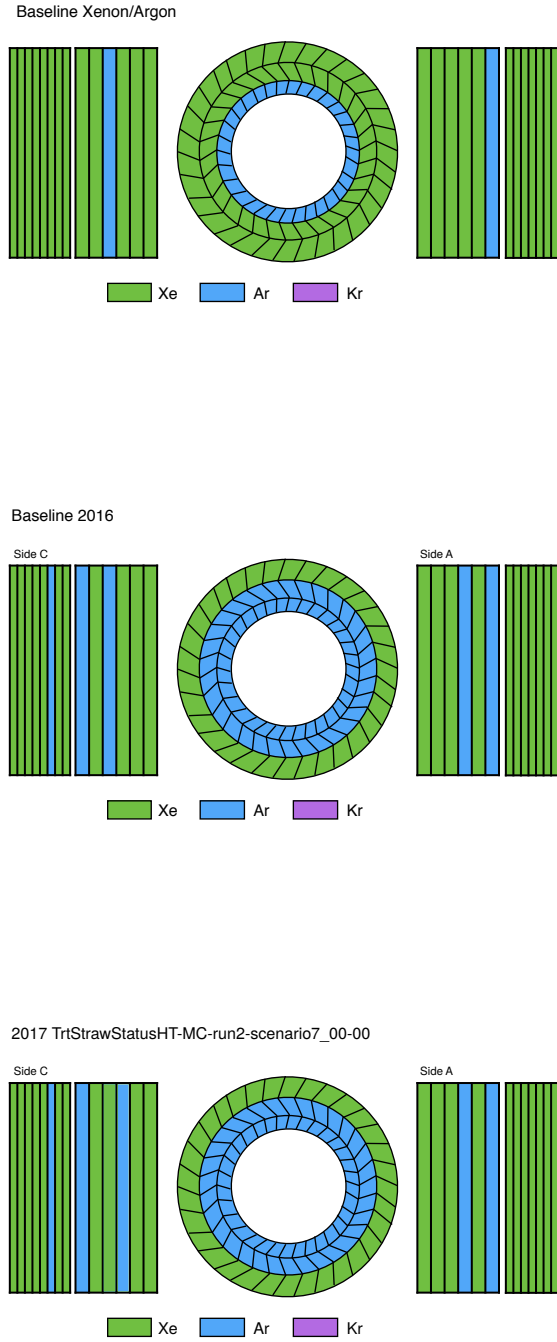


Figure 6.2: The TRT baseline gas geometries in 2015 (top), 2016 (middle) and 2017 – 2018 (bottom).

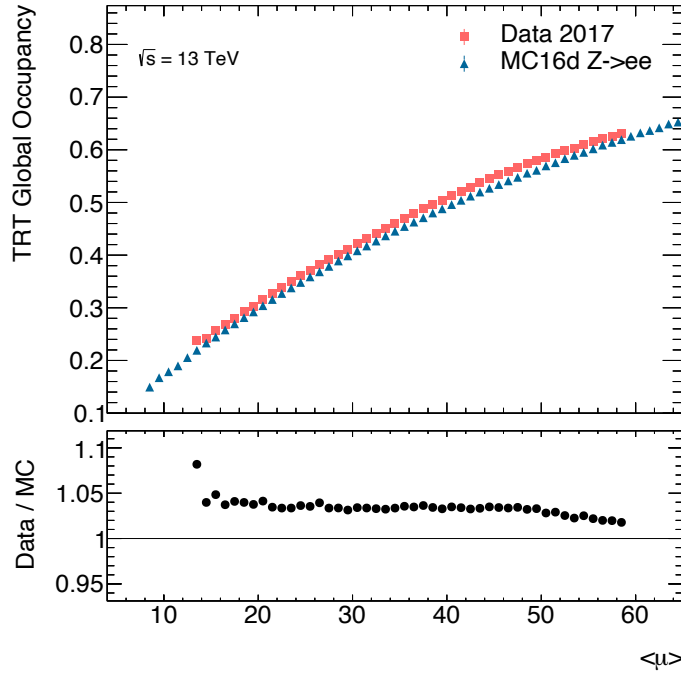


Figure 6.3: Data-MC comparison of TRT occupancy as a function of $\langle\mu\rangle$ [74].

crossing that hinders energy reconstruction and triggering, which ultimately incurs a loss of signal efficiency. The HIP kinematic regime dictates that some HIPs are slow moving such that they may experience significant delays before they arrive at the calorimeter. There are two practical consequences that follow from the delayed arrival.

Foremost, the L1 EM trigger is expected to lose efficiency for energy clusters delayed by more than 10 ns due to an increasing probability to trigger in a neighbouring collision bunch crossing crossing. The monopoles and the HECOs lose the majority of their kinetic energies during different phases of their trajectories, as modelled by the simulation (see Figure 4.2). To evaluate the trigger efficiency loss, one simplifies the velocity change by assuming the HIPs travel at constant speeds through the Inner Detector (ID). A delay of less than 10 ns corresponds to a constant HIP velocity $\beta > 0.37$ through a trajectory perpendicular to the beam line. The L1 efficiency and its dependence on the HIP truth velocity is shown in Figure 6.4. Due to geometry effects, the efficiency has a stronger dependence on the

transverse velocity $\beta \sin \theta$ than on the velocity β . A cut of $\beta \sin \theta > 0.37$ is thus applied at the analysis level to recompute the efficiency and determine the systematic uncertainty associated with late signals. On average, the systematic uncertainty is $(2 \pm 0)\%$.

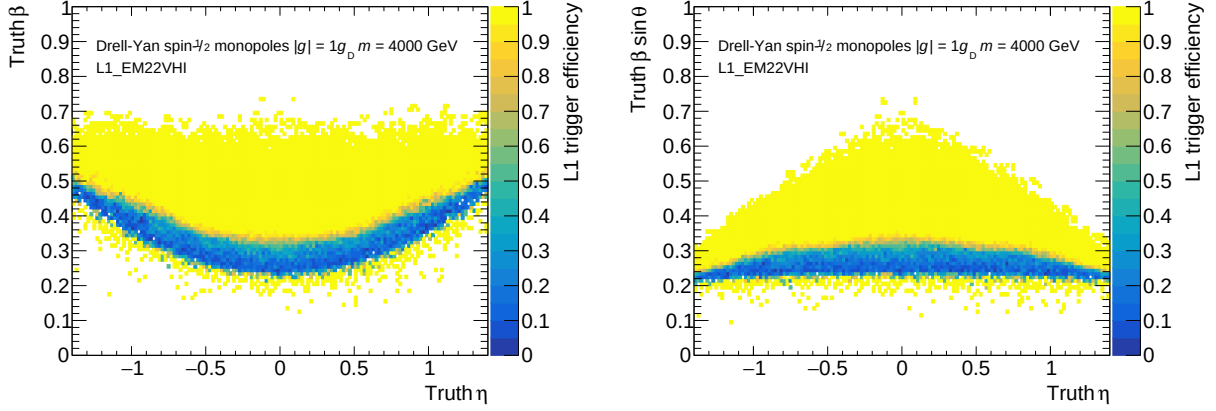


Figure 6.4: L1 efficiency as a function of truth η and truth β (left) or truth $\beta \sin \theta$ (right). Drell-Yan spin- $\frac{1}{2}$ monopoles of charge $|g| = 1g_D$ and mass 4000 GeV are used.

In addition to triggering, late signals affect the reconstruction of energy deposited in the LAr EM Calorimeter, which is optimized for particles travelling approximately at the speed of light. The reconstruction efficiency, dependent on the time delay, disagrees between data and simulation. Calorimeter experts [201] report that the systematic uncertainty in the calorimeter cluster reconstruction efficiency predicted by the Monte-Carlo simulation is no more than 2% for delays in the range of 10 – 13 ns, and less than 1% for shorter (5 – 10 ns) and longer (15 – 25 ns) time delays. To correct for the mismodelled reconstruction efficiency in the simulation, the CaloCalTopoCluster transverse energy for each simulated event is reduced by 2% at the analysis level. This systematic variation of the cluster transverse energy, introduced solely to evaluate the alternative signal selection efficiency, ultimately allows me to compute the systematic uncertainties, which turns out to be negligible.

6.2.7 HIP Corrections to Birks' Law

When charged particles cross the LAr gaps in the LAr EM Calorimeter, they ionize the active medium, which dissociates into electrons and ions that drift in the electric field toward the electrodes. The signal readout is proportional to this initial ionization that may have undergone charge recombination (see Section 4.2.4). The Birks' Law [183, 184] (see Section 4.2.4), which follows from the ICARUS measurement [185] and gets tuned to minimally ionizing particles by ATLAS, takes into account these two factors to infer the energy losses of the incident particles. However, the charge combination effect is stronger for particles with high dE/dx , hence, a different set of parameters has been derived and implemented in the ATLAS simulation to estimate the HIP energy loss in the LAr EM Calorimeter (see Section 4.2.4). These parameters are subject to up and down variations, which underlay variations of the HIP efficiencies. In order to quantify the variations in efficiencies, I requested particle-gun signal MC samples to be produced with these varying parameters. Some samples show no efficiency variation from the parameter variation. To be conservative, I quote for the systematic uncertainty the maximum among 5%, the absolute value of the systematic uncertainty from the down variation and the absolute value of the systematic uncertainty from the up variation. On average, the systematic uncertainty is $(8 \pm 0)\%$.

6.2.8 Calorimeter Crosstalk

The last source of uncertainty in the calorimeter is the simulation of energy dispersion in the clusters. Any electronic signal disturbance between neighbouring cells within the clusters in the presence of electromagnetic fields, known as crosstalk, would bias the cluster width parametrized by the w variable. Each cell is made of multiple readout electrodes that in principle also interfere with each other. Therefore, the crosstalk can occur both within

a single cell and among multiple cells, giving rise to crosstalk in varied forms: capacitive, resistive and inductive, whose effects have been studied in detail [202].

These sources of crosstalk impact different layers of the calorimeter, as summarized in Table 6.1. Due to capacitive couplings between the electrodes, the capacitive crosstalk, albeit present in all layers, is significantly enhanced in the first calorimeter layer (EM1) along the η direction where the highest segmentation is achieved. The amount of crosstalk is 7.2% and 6.3%, in barrel and end-caps, respectively, making the capacitive crosstalk particularly important for EM1. The resistive crosstalk arises as different calorimeter layers are connected by resistors mounted on the high-voltage layers of the electrodes. Since the calorimeter granularity decreases from the first to the third layer (EM3), each electrode in the second layer (EM2) is matched to more EM1 electrodes than EM3 electrodes. As a result, the resistive crosstalk is dominant between EM1 and EM2, on the order of 0.1%. The inductive crosstalk, caused by mutual inductance between the cell electrodes, turns out to overshadow the capacitive crosstalk in EM2, EM3, and between the two layers, where next-in- η neighbouring cells experience an impact on the order of 1%, while next-in- ϕ neighbouring cells show a nearly doubled impact, on the order of 1.8%.

Type	Region	Direction	Values (in %)
Capacitive	EM1-to-EM1	η	7.2 (barrel), 6.3 (end-caps)
Resistive	EM2-to-EM1	η	0.1
Inductive	EM2-to-EM2	η	1.1 (barrel), 0.6 (end-caps)
Inductive	EM2-to-EM3	η	0.9 (barrel), 1.2 (end-caps)
Inductive	EM2-to-EM2	ϕ	1.8

Table 6.1: Three types of crosstalk and their impacts [202].

The large HIP energy deposition is particularly sensitive to any crosstalk effect that may not be accurately described by the Athena simulation. Despite the next-in- η crosstalk in EM1 in the LArStripsCrossTalkCorrector class (see package LArCalorimeter/LArCalibUtils), η

crosstalk in the other layers and next-in- ϕ crosstalk are not implemented. To better describe the energy distribution among the clusters, a 1.8% energy correction is applied to next-in- ϕ neighbouring cells at the analysis level. The cluster width w is recomputed to compute the systematic efficiency. The reshuffled cell energies slightly shift the w distribution (see Figure 6.5) and ultimately degrade the efficiency if the w values of the signal events drop below the w cut. As the w cut is sufficiently large such that the signal efficiency is robust against the crosstalk effect, the uncertainty is found to be on the per-mille level, and is therefore neglected.

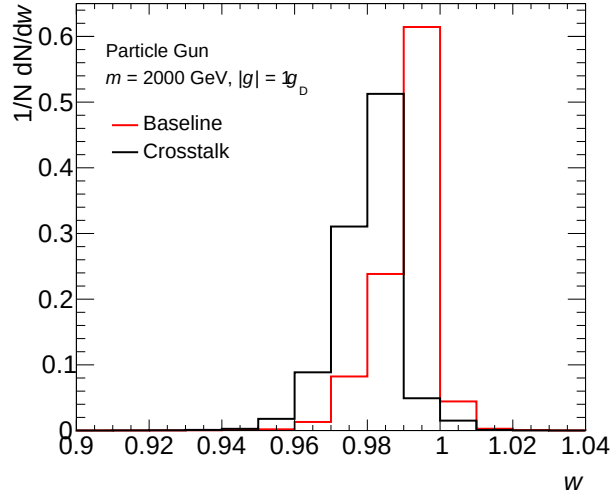


Figure 6.5: Distributions of the w variable before and after the 1.8% energy correction to the next-in- ϕ neighbouring cells.

6.2.9 Extrapolation Method

In addition to detector simulations, one contributor to the total systematic uncertainties is the extrapolation technique in estimating the efficiencies of the generator-level-only signal models, i.e., spin-0 Drell-Yan, spin-0 photon fusion and spin- $\frac{1}{2}$ photon fusion. Efficiency maps of truth transverse kinetic energy and truth pseudorapidity are produced from the fully

simulated particle-gun samples. These maps are subsequently used to sample the efficiencies of the generator-level models based on their truth kinematic properties. A truth particle is accepted if either particle in the pair passes all selection criteria.

The systematic uncertainty is computed using Drell-Yan spin- $\frac{1}{2}$ HIPs as the relative difference between the nominal efficiency from the fully simulated model and the extrapolated efficiency. It is found to be $(1 \pm 0)\%$ on average.

6.2.10 Pileup Reweighting

Last but not least, the HIP selection efficiency is sensitive to pileup that dilutes the fraction of TRT high-threshold hits in trigger and offline selection. The simulation implements a pileup distribution that is generally different from the one in data. As described in Section 5.2.2, a reweighting procedure mitigates the pileup mismodelling through weights, namely scale factors that adjust the estimated to the true hadronic activities in the detector. These weights are subject to up and down variations, which account for uncertainties of the true hadronic activities due to up and down uncertainties of the cross section of inelastic proton-proton collisions [203]. When applied to the MC events, these weight variations yield alternative pileup distributions (see Figure 6.6) and HIP selection efficiencies that allow the computation of the associated systematic uncertainties. The variation of the pileup weights result in an average systematic uncertainty of $(2 \pm 0)\%$.

6.2.11 Summary

Table 6.2 summarizes all systematic uncertainties on the signal efficiencies and their corresponding values.

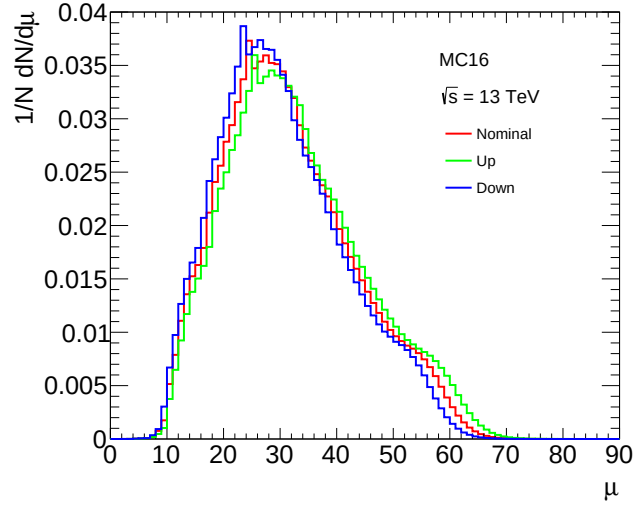


Figure 6.6: Simulated pileup μ distributions weighted by different event weights.

Systematic Uncertainty	Value in %
Detector material	9 ± 4
HIP correction to Birks' Law	8 ± 0
δ -ray production modelling	2 ± 0
TRT occupancy modelling	2 ± 0
Calorimeter arrival time	2 ± 0
Pileup correction	2 ± 0
Extrapolation	1 ± 0

Table 6.2: Summary of the systematic uncertainties on the signal efficiencies, averaged over all models, charges and masses.

6.3 Background Estimation

The ABCD method, described in Section 5.4, exploits the final selection variables f_{HT} (fraction of TRT high-threshold hits) and w (calorimeter cluster energy dispersion) to estimate the residual background after event preselection. Figure 6.7 is the ABCD plane for the 2015 – 2018 data. Two solid lines, corresponding to the final selection cuts, slice the plane into signal region A and control regions B, C and D, where

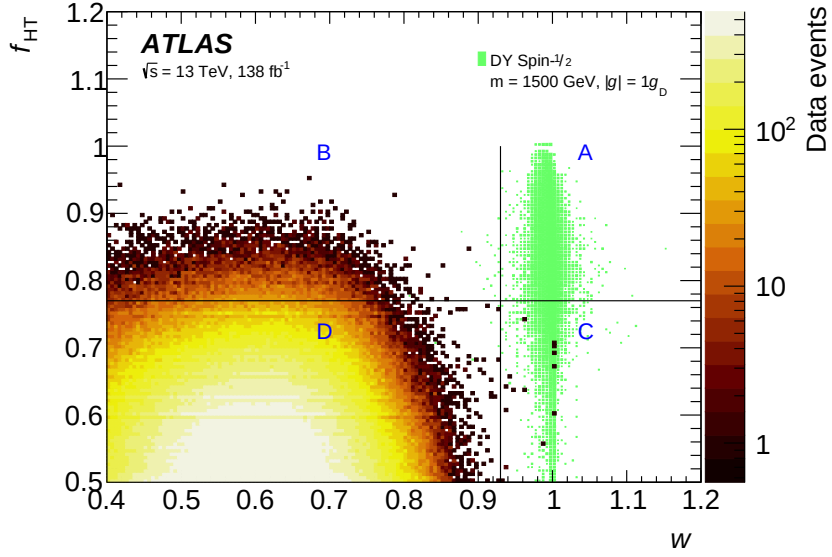


Figure 6.7: The ABCD method in the $f_{\text{HT}}-w$ plane after preselection of 2015 – 2018 data. The final selection cuts defines signal region A and control regions B, C, D. The control regions B and D are subdivided into validation regions B', B'', D', D'' [204].

- Region A (signal): $w \geq 0.93$ and $f_{\text{HT}} \geq 0.77$,
- Region B: $0.4 \leq w < 0.93$ and $f_{\text{HT}} \geq 0.77$,
- Region C: $w \geq 0.93$ and $0.5 \leq f_{\text{HT}} < 0.77$,
- Region D: $0.4 \leq w < 0.93$ and $0.5 \leq f_{\text{HT}} < 0.77$.

One dashed line further dissects the control regions B and D into validation regions B', D', B'' and D'', where the region B'' resembles the signal region A the most by construction. That is,

- Region B': $0.4 \leq w < 0.665$ and $f_{\text{HT}} \geq 0.77$,
- Region B'' (signal-like): $0.665 \leq w < 0.93$ and $f_{\text{HT}} \geq 0.77$,
- Region D': $0.4 \leq w < 0.665$ and $0.5 \leq f_{\text{HT}} < 0.77$,

- Region D'': $0.665 \leq w < 0.93$ and $0.5 \leq f_{\text{HT}} < 0.77$.

With the techniques developed earlier, three background estimates will be provided, in ascending order of statistical rigor and accuracy.

6.3.1 Background Estimate by Simple Ratio

The simplest method to estimate the background is by a simple ratio. For the ABCD plane in Figure 6.7, whose candidate events counts are tabulated in Table 6.3, the estimate in region A can be computed as

$$N_A^{\text{est}} = \frac{N_B}{N_D} \cdot N_C = 0.15 \pm 0.04, \quad (6.4)$$

where the uncertainty is incurred by statistical uncertainties of the event counts N_B , N_C and N_D , in regions B, C and D, respectively. Each statistical uncertainty is the square root of the corresponding count, as follows from Poisson statistics. The low statistics in region C dictates that the statistical uncertainty of N_C dominates the background estimate in region A.

N_A^{obs}	N_B	N_C	N_D
0	9847	15	986500

Table 6.3: Number of candidate events in data observed in regions A, B, C, D.

In order to validate the estimate, the simple-ratio method is applied to the validation regions, to yield an estimate of $N_{D''} \cdot N_{B'}/N_{D'} = 2368 \pm 28$ for the validation region B''. While this differs from the observed B'' yield in data of 2163 ± 47 by more than one standard deviation, the values are sufficiently consistent given the rudimentary nature of this test.

One caveat of this technique is that its accuracy is limited to rather uniform distributions

of the transfer factor, i.e., N_B/N_D . In the case of a non-uniform transfer factor, as discussed in Section 5.4.1, a reparametrization of the transfer factor is advisable. It turns out that the non-uniform transfer factor contributes to the largest uncertainty of the background estimate, as computed in Section 6.3.2.

6.3.2 Background Estimate by Mean Transfer Factor

One way to reparametrize the transfer factor is to consider a mean value after fitting the transfer factor as a function of w to a constant. I collaborated on this study to perform the constant fit and to extract the fit parameters. This method was first validated on the validation regions for three $|\eta|$ subranges, where the statistical uncertainty is the error of the fit, and the systematic uncertainty is the root-mean-square error with respect to the fit. As shown in Table 6.4, the estimate and the actual count in region B'' agree within uncertainties.

$ \eta $	$N_{D''} \pm (\text{stat.})$	Mean transfer factor $\pm (\text{stat.})$ $\pm (\text{sys.})$	$N_{B''} \pm (\text{stat.})$	Background estimate in B'' $\pm (\text{stat.})$ $\pm (\text{sys.})$
0 – 0.32	1258 ± 35	0.0094 ± 0.0017 ± 0.0028	14 ± 4	$12 \pm 2 \pm 4$
0.32 – 0.82	111166 ± 33	0.0012 ± 0.0001 ± 0.0001	147 ± 12	$133 \pm 6 \pm 11$
0.82 – 1.375	120036 ± 346	0.0157 ± 0.0002 ± 0.0012	2002 ± 45	$1885 \pm 21 \pm 144$
Total	-	-	2163 ± 47	$2030 \pm 22 \pm 145$

Table 6.4: Event count in validation region B'' ($N_{B''}$), and estimated background event in region B'' computed from event count in validation region D'' ($N_{D''}$) and mean transfer factor ($N_{B'}/N_{D'}$), for different $|\eta|$ subranges and in total.

Subsequently, this method was applied to the entire ABCD plane, to compute the mean

transfer factor and background estimate in region A for every $|\eta|$ slice, as tabulated in Table 6.5. The background estimate in region A is the product between the mean transfer factor and the number of candidate events in region C. The number of estimated background events, summed over all three $|\eta|$ subranges, is

$$N_A^{\text{est}} = 0.19 \pm 0.05(\text{stat.}) \pm 0.06(\text{sys.}). \quad (6.5)$$

The systematic uncertainty of 0.06 translates to a 32% uncertainty of the background estimate of 0.19, as a result of the transfer-factor non-uniformity.

$ \eta $	$N_C \pm (\text{stat.})$	Mean transfer factor $\pm (\text{stat.}) \pm (\text{sys.})$	Background estimate in A $\pm (\text{stat.}) \pm (\text{sys.})$
0 – 0.32	2 ± 1.4	0.0087 ± 0.0014 ± 0.0066	0.0174 ± 0.0126 ± 0.0180
0.32 – 0.82	2 ± 1.4	0.0013 ± 0.0001 ± 0.0002	0.0026 ± 0.0018 ± 0.0018
0.82 – 1.375	11.0 ± 3.3	0.0158 ± 0.0001 ± 0.0027	0.1738 ± 0.0524 ± 0.0602
Total	15	-	$0.19 \pm 0.05 \pm 0.06$

Table 6.5: Background estimate in region A computed from event count in region C (N_C) and mean transfer factor (N_B/N_D), for different $|\eta|$ subranges and in total.

6.3.3 Background Estimate by Maximum Likelihood Fit

One of the main motivations behind performing a maximum likelihood fit is to adequately model signal distributions in all four regions and ultimately improve the accuracy of the background estimate. When normalized by the amount of signals in the signal region A, the amount of signals in the control regions B, C, D gives rise to the signal leakage, as defined by Equation 5.8 earlier.

As shown in Figure 6.8, the signal leakage is readily computed for the fully-simulated Drell-Yan spin- $\frac{1}{2}$ model. The leakage into region C is a manifestation of the signal-selection power of the f_{HT} variable, which is obstructed by the presence of pileup and thus deteriorates, especially for low-charge HIPs. The leakage into regions B and D is less significant, except for HECOs with charge $|z| = 100$ and mass $m = 4000$ GeV. These two control regions are characterized by relatively low w values, which are expected for high-charge and high-mass HIPs that show lower efficiencies due to stopping before traversing the entire calorimeter.

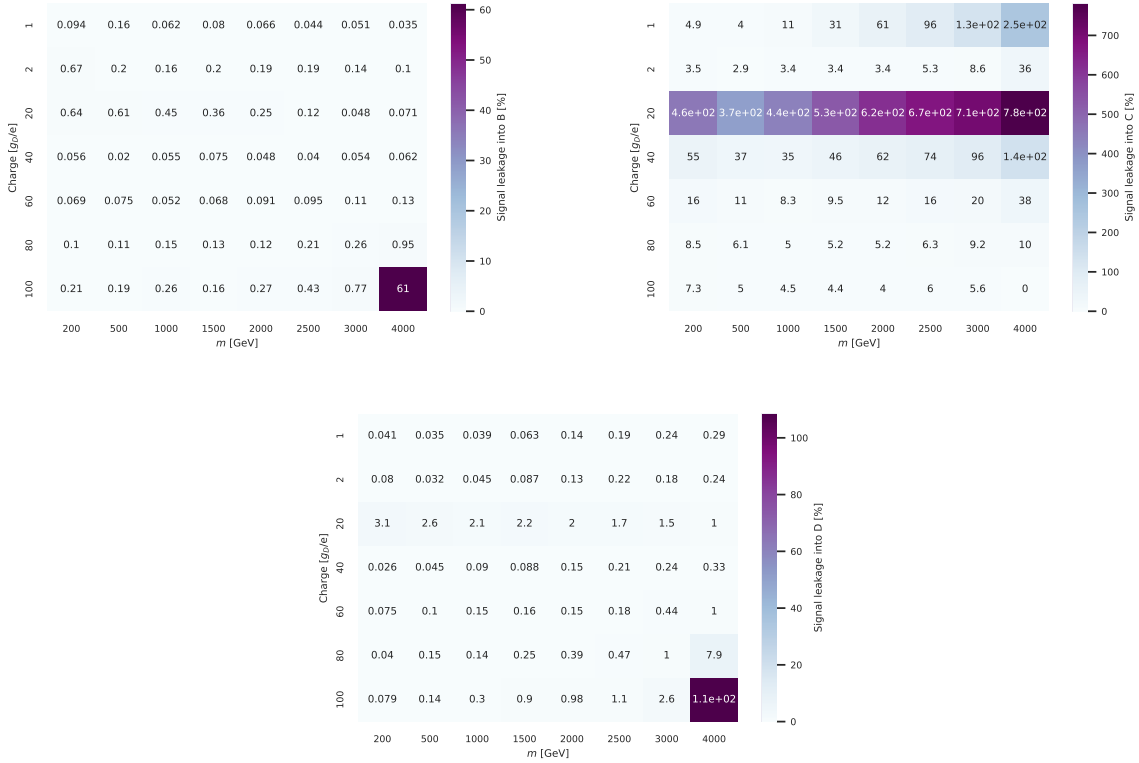


Figure 6.8: Drell-Yan spin- $\frac{1}{2}$ signal leakage into regions B, C, and D.

However, the signal leakage of the alternative models is less trivial to compute. It involves the creation of efficiency maps from particle-gun samples for all four ABCD quadrants, from which one applies the extrapolation method to sample signal yields, i.e., the fraction of MC events in a given quadrant. Expected from Figure 6.8 and confirmed with particle-

gun samples, the efficiency maps (for example see Figure 6.9) of regions B and D lack the necessary amount of statistics needed for the extrapolation method. Since the signal leakage cannot be computed for the alternative models without full simulation, they are set to zero in the likelihood fit. Two simultaneous fits are performed, both applying signal leakage in

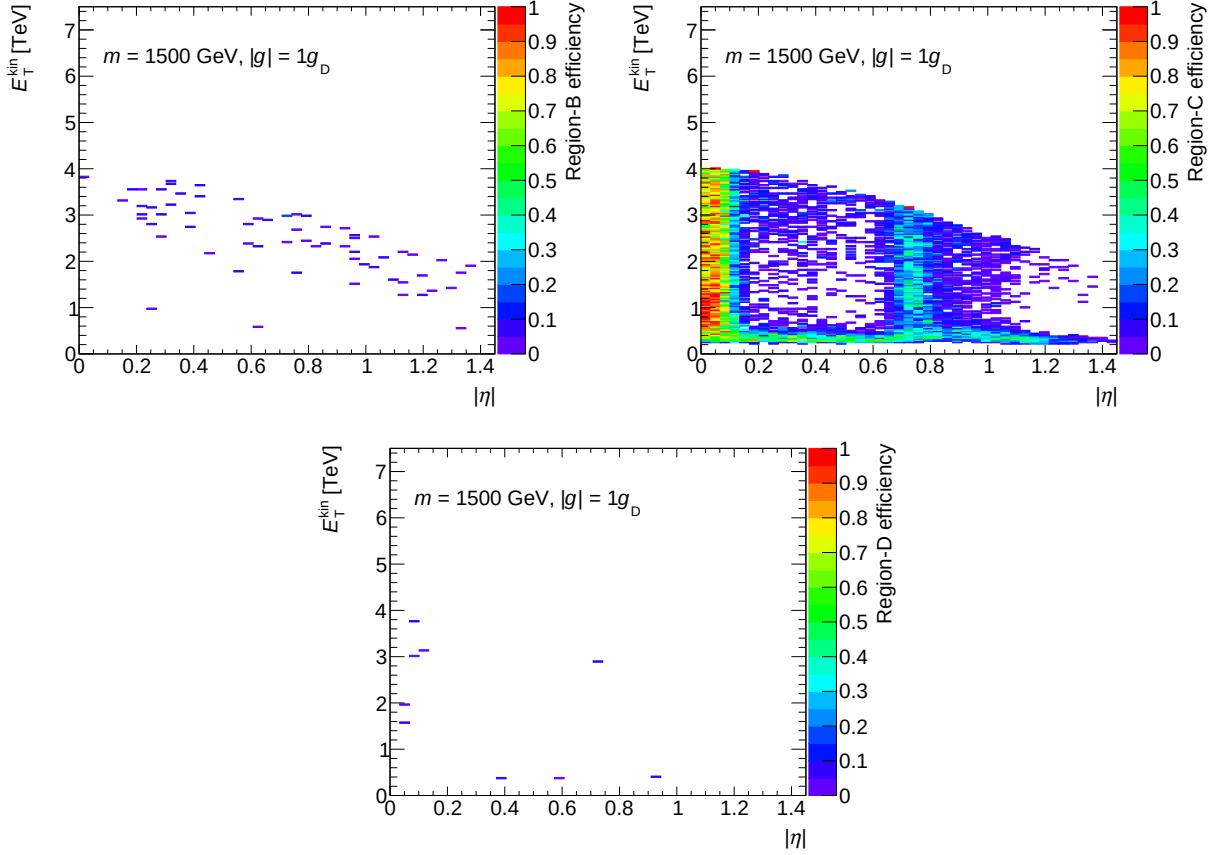


Figure 6.9: Efficiency maps of region B (top left), region C (top right) and region D (bottom) for particle-gun monopoles with charge $|g| = 1g_D$ and mass $m = 1500$ GeV. The efficiency map of region A is shown in Figure 6.1.

the maximum-likelihood fit discussed in Section 5.4.2 with an initial value of the background estimate of Equation 6.5. The first fit tests the stability of the background estimate in the presence of signal leakage. This is done without $|\eta|$ slicing through injection of different numbers of signal events in region A. The estimate turns out to agree with the pre-fit result, i.e., Equation 6.4, within uncertainties. The 32% systematic uncertainty due to transfer-

factor non-uniformity (see Equation 6.5) is neglected in the fit, thus it is assigned as a systematic uncertainty to the background estimate, resulting in an estimate of

$$N_A^{\text{est}} = 0.15 \pm 0.04(\text{stat.}) \pm 0.05(\text{sys}). \quad (6.6)$$

without $|\eta|$ slicing. The second fit is run on $|\eta|$ sliced signal and control regions without signal injection. The background estimates agree with the pre-fit result, i.e., Equation 6.5, within uncertainties. Therefore, the final background estimate with three $|\eta|$ slices is

$$N_A^{\text{est}} = 0.19 \pm 0.05(\text{stat.}) \pm 0.06(\text{sys}). \quad (6.7)$$

6.4 Result of Limits

Consistent with the background estimate in Equation 6.6 and Equation 6.7, no signal event is observed after unblinding the signal region of 2015 – 2018 data. Therefore, production cross-section limits can be set on the signal models, as introduced in Section 5.5. For each production model of HIPs of a given mass and charge, a cross-section upper limit can be computed following Equation 6.3 as

$$\sigma_{\text{UL},95}^{\text{obs}} = \frac{\mu_{\text{UL},95}^{\text{obs}} \cdot N}{\epsilon \cdot \mathcal{L}}, \quad (6.8)$$

where ϵ is the signal efficiency with statistical uncertainty (see Section 6.1) and $\mathcal{L}$ is the integrated luminosity of 138 fb^{-1} with 0.83% relative uncertainty [157]. Multiplied by the number of candidate events N predicted by simulation, the quantity $\mu_{\text{UL},95}^{\text{obs}}$, constrained by a set of systematic uncertainties on the signal efficiency (see Section 6.2) and obtained from Equation 5.15, is the upper limit of μ with higher values excluded at 95% confidence level. The derivation of $\mu_{\text{UL},95}^{\text{obs}}$ relies on multiple datasets that are collected by *pseudo-experiments*.

With the true number of background events ν_t being determined, each pseudo-experiment randomly samples the number of background events ν_0 from a Gaussian distribution with mean ν_t and width $\delta\nu$ and the number of observed events n_0 from a Poisson distribution with mean ν_t . and subsequently computes q_μ from the generated ν_0 and n_0 . Subsequently, the parameter of interest μ is measured in each dataset from the generated ν_0 and n_0 . As the number of datasets becomes sufficiently large, the true value of μ is expected to be contained in the range of varying estimates of μ within a certain degree of probability. With 5000 pseudo-experiments, CL_s is computed as a function of μ for each model of a given HIP mass and charge. The μ value with a CL_s value of 0.05 is interpolated as the upper limit of the signal strength $\mu_{\text{UL},95}^{\text{obs}}$. Implemented via the HISTFITTER framework [195], this method yields cross-section limits at 95% confidence level as a function of HIP mass for the different production models and charges, as shown in Figure 6.10 and Figure 6.11. These limits are obtained without $|\eta|$ slicing the ABCD plane, as the limits computed with and without $|\eta|$ -slicing for Drell-Yan spin- $\frac{1}{2}$ HIPs only differ by 5.5% on average. Except for the Drell-Yan spin- $\frac{1}{2}$ HIPs whose signal leakage in the control regions is considered, the other production models only consider signal-region information. This is not expected to bring significant bias to the limits of the alternative models, as the limits computed with and without the signal leakage for the Drell-Yan spin- $\frac{1}{2}$ HIPs differ by less than 1%.

Lower mass limits are as tabulated in Table 6.6, such that masses below these limits are excluded at 95% confidence level. In most cases, the mass limits correspond to intersections between the cross-section upper limits and the expected cross sections. However, if the cross section upper limit curve does not intersect the expected cross section curve, the highest mass for which a cross section upper limit was computed is interpreted as the mass limit.

As shown in Figure 6.12, the Drell-Yan limits supersede the early Run 2 result, thanks to the much greater luminosity of the full Run 2 data sample, i.e., 138 fb^{-1} in comparison to the 2015 – 2016 data sample of 34.4 fb^{-1} . The photon-fusion mechanism shows stronger

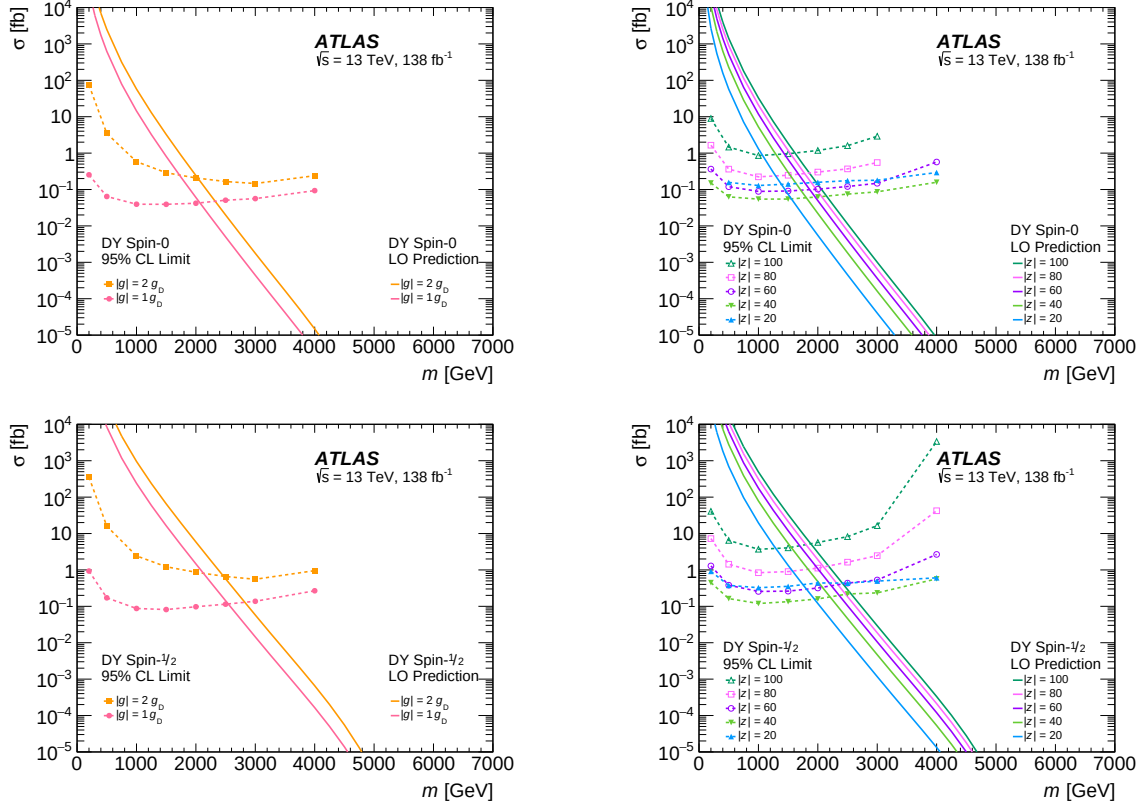


Figure 6.10: Upper limits on production cross section at 95% confidence level for Drell-Yan (DY) spin-0 and spin- $\frac{1}{2}$ HIPs as a function of mass (dashed lines with markers), overlaid with theoretical cross sections (solid lines). Mass points for which the selection efficiency is below 0.1% are not shown [204].

	HIP lower mass limits [GeV]						
	$ g = 1g_D$	$ g = 2g_D$	$ z = 20$	$ z = 40$	$ z = 60$	$ z = 80$	$ z = 100$
DY spin-0	2097	2055	1425	1833	1901	1791	1651
DY spin- $\frac{1}{2}$	2554	2482	1801	2216	2239	2115	1911
PF spin-0	3350	3469	2143	2828	2939	2801	2500
PF spin- $\frac{1}{2}$	3613	3737	2480	3117	3142	3000	2500

Table 6.6: Lower mass limits (in GeV) at 95% confidence level in models of spin-0 and spin- $\frac{1}{2}$ Drell-Yan (DY) and photon-fusion (PF) pair production of HIPs [204].

exclusion power than the Drell-Yan mechanism, since the theoretical cross sections of photon fusion are higher than those of the Drell-Yan mechanism. In sight of higher centre-of-mass

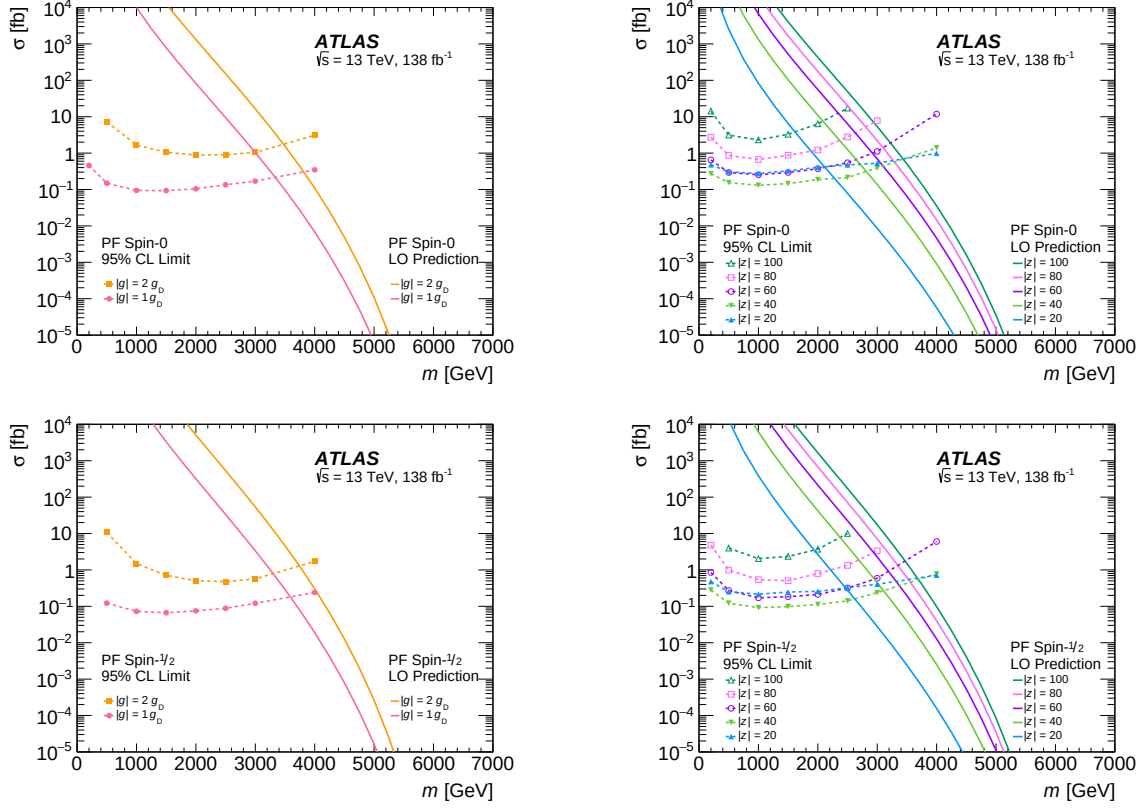


Figure 6.11: Upper limits on production cross section at 95% confidence level for photon-fusion (PF) spin-0 and spin- $\frac{1}{2}$ HIPs as a function of mass (dashed lines with markers), overlaid with theoretical cross sections (solid lines). Mass points for which the selection efficiency is below 0.1% are not shown [204].

energy and total luminosity in Run 3, the photon-fusion mechanism is expected to exclude more parameter space, if its selection efficiency does not drop significantly in the presence of higher pileup. Therefore, how to improve the selection efficiency will be a subject of future work. Another direction with which one may expand the current work is to consider the production of HIPs of mass less than 200 GeV. Although the current search is not sensitive to this mass range, I will show the phenomenological prospect of exploiting a different collision system in the subsequent chapter.

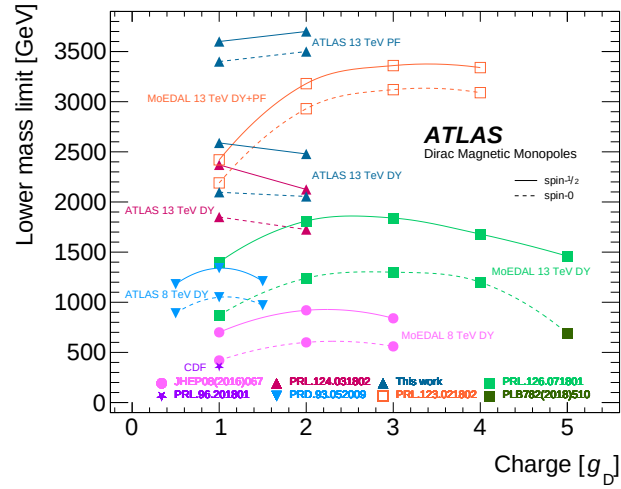


Figure 6.12: Comparison of monopole mass limits by LHC searches in Run 1 and Run 2 proton-proton collisions assuming the Drell-Yan (DY) production mechanism. A Tevatron measurement by CDF in proton-antiproton collisions is also shown. [204]

Chapter 7

Highly Ionizing Particle Search in Heavy Ion Collisions

This chapter explores heavy-ion collisions as a possible extension for ATLAS Run 3 highly ionizing particle (HIP) searches. The heavy-ion collisions are of great interest due to the large electromagnetic fields associated with the high-charge ions moving at high velocities. Ultra-peripheral collisions (UPC) could produce magnetic monopoles via the photon-fusion (PF) mechanism. Furthermore, UPCs have a characteristic signature with low particle multiplicities, making them easy to select and analyze. Following the development of the photon-fusion mechanism in proton-proton collisions (see Chapter 2), I proceeded to model the production of HIPs in UPC events by modelling the photons emitted by the heavy ions. With the model set up in MADGRAPH5_aMC@NLO [147], I calculated the HIP production cross sections and kinematic properties at the LHC design centre-of-mass collision energy. This material has been published in [141].

This chapter is outlined as follows. Section 7.1 introduces heavy-ion collisions, explains ultraperipheral collisions as one collision type in such collisions, and compares the parameters of collisions of various heavy ions to those of protons. Section 7.2 discusses the model used

to simulate HIP production in heavy-ion collisions, followed by cross-section results and representative kinematic distributions. Section 7.3 contemplates the pros and cons of such a search and how to tailor the search strategy to meet the unique challenges.

7.1 Ultraperipheral Heavy-Ion Collisions

The LHC not only collides protons but also heavy ions. Heavy-ion collisions have many distinctive properties that make them a unique environment for studying a wide range of interactions, including the production of magnetic monopoles. The effective interaction in heavy-ion collisions is characterized by the impact parameter of the colliding ions. When the ions interact with impact parameters less than or equal to twice the nuclear radii, the strong force governs the interaction between them. These are known as central collisions and their extraordinarily high temperatures can produce a quark-gluon plasma. However, when the two ions pass each other with an impact parameter much larger than twice their radii, strong interactions between the nucleons are no longer possible. Nevertheless, even at large impact parameters electromagnetic processes between the ions can occur. In these so-called ultraperipheral collisions (UPC), photons radiated from the ions interact via the photon-fusion process while the ions remain intact.

The prospect of monopole production in heavy-ion collisions has a long history of interest. Initially, it was envisioned to occur via thermal processes of central collisions where the massive monopoles would couple to the quark-gluon plasma [205] via photons. Later, monopole production was postulated in ultraperipheral elastic photon-fusion processes in electron-electron, proton-proton and lead-lead collisions [143], motivated by the enhanced cross sections due to the high electric charges of the heavy ions over protons. Recently, monopole production has been considered again in ultraperipheral heavy-ion collisions [90–96], where the strong magnetic field from the ultrarelativistic high-charge ions passing each

other may produce monopoles via the magnetic-dual mechanism of the thermal Schwinger process [206]. This production mechanism has been considered in an experimental search for the first time by the MoEDAL Collaboration [81]. The present study first focuses on the electromagnetic radiation emitted by the ultrarelativistic ions, as in [143], then continues to formulate a model of monopole and HECO production in (the magnetic dual of) Quantum Electrodynamics, assuming the effective interaction is photon fusion.

7.1.1 Electromagnetic Radiation Emitted by Relativistic Heavy Ions

The *equivalent photon method* describes the electromagnetic radiation emitted in heavy-ion collisions. The original formalism, proposed by Enrico Fermi in 1924 [207], treated the time-varying electromagnetic field as a flux of virtual photons. Later, it was adapted by Weizsäcker and Williams [208, 209] for the photon fluxes of relativistic ions.

One considers a particle with electric charge Z as approximately a point particle. At ultrarelativistic velocities, it produces electric and magnetic fields ($\mathbf{E}$ and $\mathbf{B}$) such that the total energy the fields carry propagates as a flux of equivalent photons with the flux direction and magnitude dictated by $\mathbf{E} \times \mathbf{B}$. The resulting photon density spectrum dn_γ possesses a Z^2 dependence and can be defined for a given pair of energy ω and impact parameter b in the plane transverse to its direction of motion. The number density over the entire transverse plane is obtained by integrating over all possible impact parameters b . Therefore, the number density per energy is

$$\frac{dn_\gamma}{d\omega}(\omega) = 2\pi \int_{b_{\min}}^{\infty} \frac{dn_\gamma}{d\omega}(\omega, b) b db. \quad (7.1)$$

Defined with respect to one particle, this impact parameter, denoted as b hereafter, should

be distinguished from the collision impact parameter that categorizes the collision type. If the two colliding particles have the same size, the particle impact parameter b is half of the collision impact parameter on average.

Now one considers the scenario of two charged particles moving towards each other at ultrarelativistic velocities. In an ultraperipheral collision, the nuclei remain intact and the interactions take place at a finite distance, which in turn restricts the kinematics of the virtual photons due to the uncertainty principle. Consequently, the emitted photon's wavelength given by the impact parameter b has a lower bound of the nuclear radius R_A , which corresponds to a maximal transverse momentum of $\hbar/R_A$. Since the photon's off-shell mass is suppressed by the impact parameter, its virtuality decreases with the size of the charged particle. In the longitudinal direction, the photon's momentum is boosted by a Lorentz factor, reaching an effective maximum of $\gamma\hbar/R_A$. It follows from the $1/R_A$ dependence that the photon energy bound is more stringent in heavy-ion collisions than in proton-proton collisions.

7.1.2 Collisions at the LHC

Although theories describing the production of magnetic monopoles and HECOs provide little guidance as to their masses, collision energy plays an enormous role in the mass range of interest when it comes to collider searches. In ultraperipheral heavy-ion collisions at the LHC, the accessible mass range is restricted by the maximum energy of the emitted photon, $\gamma\hbar/R_A$, which is 82 GeV for lead-lead collisions and thus is much smaller than the mass upper bound of 4000 GeV the current analysis has probed in proton-proton collisions.

Table 7.1 compares the parameters of ultraperipheral collisions between different particles at the LHC design beam energy, where the ions are fully stripped with no remaining electrons. The proton-proton collision centre-of-mass energy $\sqrt{s}$ is equal to the sum of the proton beam

energies, i.e., 13 TeV in Run 2 at the LHC. The heavy-ion centre-of-mass energy $\sqrt{s_{\text{NN}}}$ is determined by the energy carried by the protons in the colliding beams, since only the protons in the ions are accelerated. As a consequence, given the same beam energy, $\sqrt{s_{\text{NN}}}$ for heavy-ion collisions is lower than $\sqrt{s}$ for proton-proton collisions. Given the $1/R_A$ dependence of the maximal transverse momenta $p_{\text{T}}^{\text{max}}$ for the different collision systems, the p_{T} profile of the photons is highly suppressed in a heavy-ion collision compared to a proton-proton collision.

	E_{beam} [TeV]	$\sqrt{s_{\text{NN}}}$ [TeV]	γ	R_A [fm]	Z^4	Z/A	$p_{\text{T}}^{\text{max}}$ [MeV]	$p_{\text{L}}^{\text{max}}$ [GeV]	$\sqrt{s_{\gamma\gamma}^{\text{max}}}$ [GeV]
p	7.0	14	7.5×10^3	0.8	1.0	1.0	237	1.8×10^3	3.6×10^3
$^{16}_8\text{O}$	7.0	7.0	3.8×10^3	3.0	4.1×10^3	0.50	65	245	490
$^{129}_{54}\text{Xe}$	7.0	5.9	3.1×10^3	6.1	8.5×10^6	0.42	32	102	204
$^{208}_{82}\text{Pb}$	7.0	5.5	3.0×10^3	7.1	4.5×10^7	0.39	28	82	164

Table 7.1: Parameters in different ultraperipheral collision systems with the same beam energy E_{beam} : nucleon-nucleon centre-of-mass energy $\sqrt{s_{\text{NN}}}$, Lorentz factor γ , nuclear radius R_A [210], charge-induced photon luminosity enhancement Z^4 , charge-to-mass ratio Z/A , maximal transverse momentum $p_{\text{T}}^{\text{max}}$, maximal longitudinal momentum $p_{\text{L}}^{\text{max}}$ and maximal photon-photon centre-of-mass energy $\sqrt{s_{\gamma\gamma}^{\text{max}}}$.

The LHC has a long-standing interest in collisions between lighter ions, which have an advantage in ultraperipheral collisions. In Run 2, the LHC delivered xenon-xenon collisions in addition to lead-lead collisions. Different collision systems have different ratios of atomic number Z and mass number A . Hence, $\sqrt{s_{\text{NN}}}$ is higher for an ion with a higher ratio of atomic number Z to mass number A , for a fixed beam energy. Heavy nuclei with high Z values have significantly stronger electrostatic interactions between the protons extending beyond their sizes. Hence, higher fractions of neutrons are necessary for the short-range strong interaction

between nucleons to keep the ion stable. In other words, light ions such as oxygen have higher Z/A values than the commonly considered xenon and lead ions. The energy fraction carried by the protons will increase and result in higher photon-photon collisions energies. The LHC foresees the addition of oxygen-oxygen collisions to its heavy-ion collision program in Run 3.

7.2 MadGraph Implementation and Results

One key component of the MADGRAPH5_aMC@NLO simulation of ultraperipheral HIP production is the distribution of photons emitted by the ions, as it acts as the photon parton distribution function for this collision mode. I implemented and exploited the distribution of emitted photons to compute the cross sections and kinematics for spin-0 and spin- $\frac{1}{2}$ production of monopoles and HECOs. Assuming the LHC design collision energy of $\sqrt{s} = 14$ TeV, these results are compared to those for the Drell-Yan and photon-fusion mechanisms in proton-proton collisions. Since the photons emitted by the lead ion beams are kinematically bounded by 82 GeV at the lead-lead collision energy of $\sqrt{s_{\text{NN}}} = 5.5$ TeV, the HIP mass range accessible in ultraperipheral collisions is much lower than that of proton-proton collisions. Hence, $|g| = 1g_{\text{D}}$ monopoles and $|z| = 60$ HECOs with masses below 82 GeV are examined.

7.2.1 Implementation of Equivalent Photon Fluxes

In heavy-ion ultraperipheral collisions, the total cross section for HIP production factorizes into the product of the parton-level cross section convoluted with the equivalent photon spectra from the two colliding ion beams. The parton-level interaction is the same photon fusion as discussed for proton-proton collisions, while the photon spectrum is uniquely parametrized for a given nuclear charge Z that follows Equation 7.1 in the Weizsäcker-

Williams approach [211],

$$f_\gamma(x) = \frac{Z^2\alpha}{\pi} \frac{1}{x} [2x_i K_0(x_i) K_1(x_i) - x_i^2 (K_1^2(x_i) - K_0^2(x_i))], \quad (7.2)$$

where $x_i = xm_N b_{\min}$ builds in the Lorentz boost effect through the atomic mass m_N and excludes hadronic interactions via the minimal impact parameter requirement $b_{\min} = R_A$, while $K_{0,1}$ are the modified Bessel functions of the second kind of zero and first order that encode longitudinal and transverse polarizations of the spectrum. In principle, there is no constraint on the intactness of the nuclei for an ultraperipheral collision to take place, as long as there is no hadronic component in the same collision event. However, in a typical calculation, a cut of the nuclear radius is applied to the impact parameter to ensure the absence of hadronic activity. The nuclear charge dependence of the photon spectrum scales as Z^2 , enhancing the cross section in heavy-ion collisions by Z^4 compared to photon fusion in proton-proton collisions. The photons are transversely polarized for relativistic ions, as parametrized by the dominance of K_1 over K_0 .

The implemented $K_{0,1}$ factors are polynomial approximations of the true $K_{0,1}$ functions, with a precision on the order of 10^{-7} [212]. For a lead ion with 82 protons and 126 neutrons, the atomic mass m_N is set to 208 times the proton mass. The radius R_A is set to 6.64 fm.

7.2.2 Cross Sections and Kinematic Distributions

The production cross sections of spin-0 $|g| = 1g_D$ monopoles and $|z| = 60$ HECOs are depicted in Figure 7.1, where three different collision modes are presented. The lead-lead ultraperipheral production is compared to proton-proton production via the Drell-Yan and the photon-fusion mechanisms. The ultraperipheral production dominates the proton-proton production mechanisms at the considered mass range.

The ultraperipheral production cross sections of $|g| = 1g_D$ monopoles and $|z| = 60$

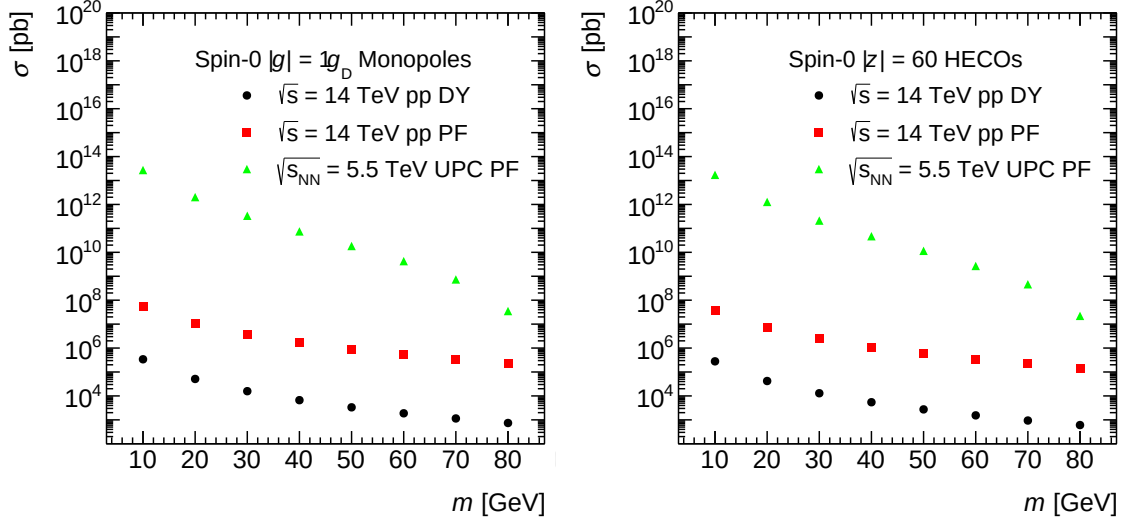


Figure 7.1: Cross sections for spin-0 HIPs via the Drell-Yan (DY) and photon-fusion (PF) mechanisms in proton-proton (pp) collisions and via photon fusion (PF) in lead-lead ultra-peripheral collisions (UPC).

HECOs are shown in Figure 7.2, where the spin- $\frac{1}{2}$ HIPs have larger cross sections than the spin-0 HIPs.

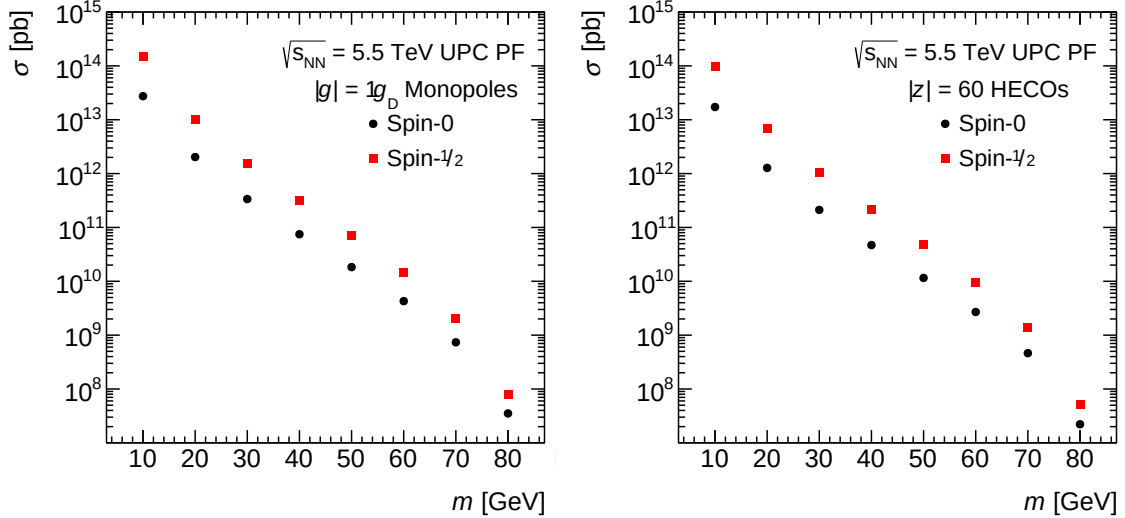


Figure 7.2: Cross sections for spin-0 and spin- $\frac{1}{2}$ HIPs via photon fusion (PF) in lead-lead ultraperipheral collisions (UPC).

The kinematic distributions for ultraperipheral production of $|g| = 1g_D$ monopoles with

mass 40 GeV are presented in Figure 7.3. Aforementioned in Chapter 2, the choice of the HIP charge does not change the kinematics. The transverse momentum distribution is softer for spin-0 HIPs than for spin- $\frac{1}{2}$ HIPs. The pseudorapidity distributions are very similar in the two spin scenarios.

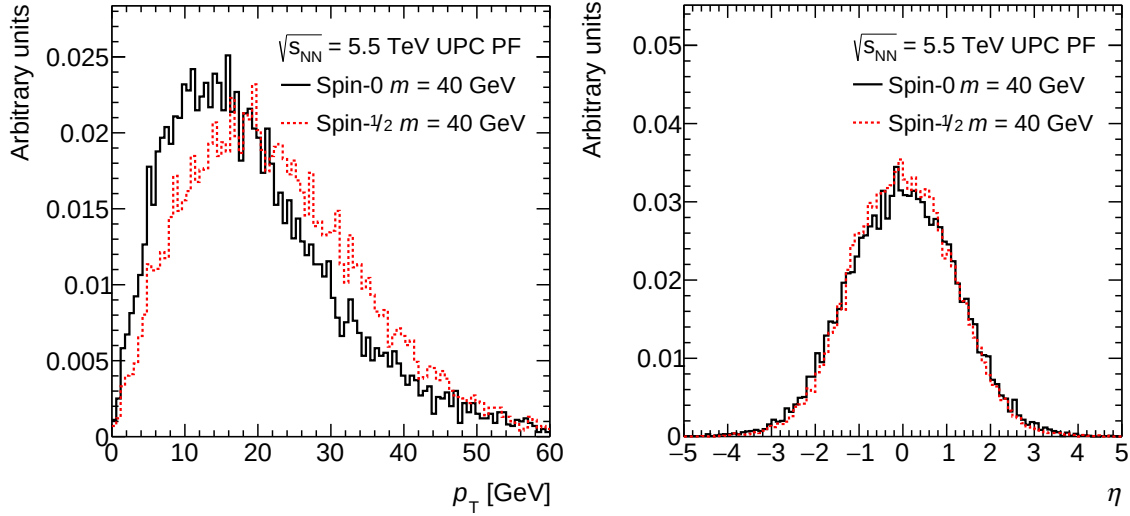


Figure 7.3: Without loss of generality, $|g| = 1g_D$ monopoles were chosen to show the transverse momentum p_T (left) and pseudorapidity η (right) distributions of spin-0 (black solid line) and spin- $\frac{1}{2}$ (red dotted line) HIPs with mass 40 GeV produced in lead-lead ultraperipheral collisions (UPC).

7.3 Summary and Outlook

The HIP kinematics are significantly more limited in the ultraperipheral collisions, hence, for comparison purposes their production in both collision types is examined in a mass range that is accessible in ultraperipheral collisions. In this mass range, their production in lead-lead ultraperipheral collision is significantly enhanced due to the large nuclear charge compared to the photon-fusion process in proton-proton collisions. Although this difference plays no role at the parton level on the matrix elements, it enhances the luminosity of the effective photon collisions. In principle, such an enhancement translates to a fast onset of

experimental sensitivity as more data is gathered in Run 3.

However, the low transverse momenta of these HIPs make it very challenging to measure them with the ATLAS Detector. In particular, it is practically impossible to conduct a search with the current analysis strategy, as according to a detector simulation I ran these HIPs will come to a stop before they arrive at the Calorimeter and even the TRT, which are extensively used for triggering and reconstruction. Undoubtedly this exotic signature calls for innovative triggers and search strategies.

Chapter 8

Conclusion

Speaking of “the safest bets one can make about physics not yet seen” [213], magnetic monopoles arise in various theoretical frameworks and models, ranging from Grand Unified Theory to Electromagnetism. If one is found, the broken electric-magnetic dual symmetry in Maxwell’s Equations will be restored. Models that predict monopole masses around the TeV scale are particularly intriguing for searches at the LHC. These monopoles, with high equivalent electric charges, continuously lose a large amount of kinetic energy before they come to a halt through ionizing surrounding material, analogous to stable high-electric-charge objects. As a result, monopoles and HECOs, known by their highly ionizing nature as highly ionizing particles, give rise to a distinctive experimental signature that allows for the selection of these particles among other particles emerging from the LHC collisions.

The successful Run 2 proton-proton collisions at 13 TeV during 2015 – 2018 culminated in a HIP data sample of 138 fb^{-1} with the ATLAS Detector at the LHC. The bulk of this work focuses on the workflow and result of a search using this sample. Since no deviations from Standard-Model predictions are found, limits on benchmark models are computed for HIPs with different charges. The Drell-Yan production model is further constrained with data collected in later Run 2, resulting in the extension of cross-section limits beyond those

obtained in the 2015 – 2016 search. Furthermore, the photon-fusion model has been considered for the first time in an ATLAS HIP search. These HIPs populate a slightly different phase space, which resulted in higher efficiency and cross-section limits than their Drell-Yan counterparts. The ATLAS sensitivity, based on an energy-loss approach, is highly dependent on the HIP charge. As far as the low-charge HIPs are concerned, the ATLAS results continue to show its unmatched constraining power on the HIP parameter space.

Testifying to the versatility of the ATLAS Experiment, the current HIP search utilizes dedicated simulation and non-standard reconstruction. The simulation distinguishes magnetic charges from electric charges and maps out their different bending directions in a magnetic field. The data, both real and simulated, is saved into a special format that allows for the analysis of ionization-dense regions of finite sizes instead of the generic particle tracks. The HIP signature, despite being simulated to the best possible accuracy, is obscured by the presence of pileup resulting from the unprecedented LHC collision rate. The selection of HIP candidates is designed to strike a balance between a simulation-driven signal efficiency and a data-driven background rejection.

As this work is being completed, the LHC has resumed its operations with Run 3, launching proton-proton collisions at higher centre-of-mass energy and pileup. The HIP analysis envisages the continuation of its role in the quest for magnetic monopoles and HECOs. It is entirely conceivable that increasing collision energy and luminosity will improve the discovery potential or the parameter constraints and thus supersede the current limits. However, the techniques and strategies developed in the Run 2 search will carry over to future studies. On the other hand, as more and more parameter regions have been excluded by experimental efforts, it may be time to think about what has been neglected.

The LHC and the ATLAS Detector, by construction, offer great discovery potential for high transverse-momentum (p_T) physics. The HIP searches, resting on anomalously large energy losses, have also been led in this direction. Although the proton-proton collisions

seem to naturally motivate just the direction we have taken, the heavy-ion collisions span a new territory with different features we ought to explore. Meaning, in order to conduct a search in ultraperipheral heavy-ion collisions, the search strategy will have to be reinvented from triggering to offline event selection to work with low- p_T HIPs. Hopefully, as further progress is made in understanding this and as the Run 3 data rolls in, it will be possible to showcase real experimental searches, rather than merely phenomenological speculation.

Appendix A

Complete Signal Efficiencies

Mechanism	Spin	Charge [g_D/e]	Mass [GeV]	Efficiency [%]	$\pm$ Error [%]
DY	0	1	200	7.94	0.06
DY	0	1	500	33.0	0.2
DY	0	1	1000	52.3	0.2
DY	0	1	1500	54.0	0.2
DY	0	1	2000	49.1	0.2
DY	0	1	2500	43.1	0.2
DY	0	1	3000	37.8	0.2
DY	0	1	4000	22.9	0.2
DY	0	2	200	0.0281	0.0002
DY	0	2	500	0.591	0.003
DY	0	2	1000	3.47	0.02
DY	0	2	1500	7.09	0.04
DY	0	2	2000	9.93	0.07

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Mechanism	Spin	Charge [g_D/e]	Mass [GeV]	Efficiency [%]	$\pm$ Error [%]
DY	0	2	2500	12.38	0.08
DY	0	2	3000	14.32	0.09
DY	0	2	4000	9.02	0.07
DY	0	20	200	6.4	0.1
DY	0	20	500	13.2	0.1
DY	0	20	1000	16.6	0.2
DY	0	20	1500	15.4	0.2
DY	0	20	2000	13.9	0.2
DY	0	20	2500	11.7	0.1
DY	0	20	3000	11.1	0.1
DY	0	20	4000	7.8	0.1
DY	0	40	200	14.1	0.1
DY	0	40	500	32.1	0.2
DY	0	40	1000	39.8	0.2
DY	0	40	1500	38.0	0.2
DY	0	40	2000	33.8	0.2
DY	0	40	2500	28.3	0.2
DY	0	40	3000	24.7	0.2
DY	0	40	4000	12.9	0.1
DY	0	60	200	5.74	0.05
DY	0	60	500	17.4	0.1
DY	0	60	1000	24.3	0.1

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Mechanism	Spin	Charge [g_D/e]	Mass [GeV]	Efficiency [%]	$\pm$ Error [%]
DY	0	60	1500	23.2	0.1
DY	0	60	2000	20.4	0.1
DY	0	60	2500	17.0	0.1
DY	0	60	3000	13.7	0.1
DY	0	60	4000	3.77	0.06
DY	0	80	200	1.213	0.007
DY	0	80	500	5.78	0.04
DY	0	80	1000	9.11	0.07
DY	0	80	1500	8.57	0.07
DY	0	80	2000	7.01	0.07
DY	0	80	2500	5.69	0.06
DY	0	80	3000	3.96	0.05
DY	0	80	4000	0.3	0.01
DY	0	100	200	0.224	0.002
DY	0	100	500	1.38	0.01
DY	0	100	1000	2.48	0.02
DY	0	100	1500	2.31	0.03
DY	0	100	2000	1.7	0.02
DY	0	100	2500	1.25	0.02
DY	0	100	3000	0.69	0.01
DY	0	100	4000	0.0026	0.0006
DY	$\frac{1}{2}$	1	200	2.3	0.01

Continued on next page

Mechanism	Spin	Charge [g_D/e]	Mass [GeV]	Efficiency [%]	$\pm$ Error [%]
DY	$\frac{1}{2}$	1	500	12.38	0.06
DY	$\frac{1}{2}$	1	1000	23.7	0.1
DY	$\frac{1}{2}$	1	1500	25.1	0.1
DY	$\frac{1}{2}$	1	2000	22.1	0.1
DY	$\frac{1}{2}$	1	2500	18.6	0.1
DY	$\frac{1}{2}$	1	3000	15.9	0.1
DY	$\frac{1}{2}$	1	4000	7.86	0.07
DY	$\frac{1}{2}$	2	200	0.00585	3×10^{-5}
DY	$\frac{1}{2}$	2	500	0.1218	0.0006
DY	$\frac{1}{2}$	2	1000	0.776	0.004
DY	$\frac{1}{2}$	2	1500	1.668	0.009
DY	$\frac{1}{2}$	2	2000	2.49	0.02
DY	$\frac{1}{2}$	2	2500	3.15	0.02
DY	$\frac{1}{2}$	2	3000	3.67	0.02
DY	$\frac{1}{2}$	2	4000	2.29	0.02
DY	$\frac{1}{2}$	20	200	2.24	0.04
DY	$\frac{1}{2}$	20	500	5.51	0.07
DY	$\frac{1}{2}$	20	1000	6.67	0.08
DY	$\frac{1}{2}$	20	1500	6.35	0.07
DY	$\frac{1}{2}$	20	2000	5.46	0.07
DY	$\frac{1}{2}$	20	2500	4.89	0.06
DY	$\frac{1}{2}$	20	3000	4.46	0.06

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Mechanism	Spin	Charge [g_D/e]	Mass [GeV]	Efficiency [%]	$\pm$ Error [%]
DY	$\frac{1}{2}$	20	4000	3.25	0.05
DY	$\frac{1}{2}$	40	200	4.68	0.04
DY	$\frac{1}{2}$	40	500	12.93	0.08
DY	$\frac{1}{2}$	40	1000	17.16	0.09
DY	$\frac{1}{2}$	40	1500	15.86	0.08
DY	$\frac{1}{2}$	40	2000	12.9	0.08
DY	$\frac{1}{2}$	40	2500	9.09	0.06
DY	$\frac{1}{2}$	40	3000	8.41	0.06
DY	$\frac{1}{2}$	40	4000	3.71	0.04
DY	$\frac{1}{2}$	60	200	1.61	0.01
DY	$\frac{1}{2}$	60	500	5.32	0.03
DY	$\frac{1}{2}$	60	1000	8.36	0.05
DY	$\frac{1}{2}$	60	1500	7.81	0.05
DY	$\frac{1}{2}$	60	2000	6.51	0.04
DY	$\frac{1}{2}$	60	2500	5.16	0.04
DY	$\frac{1}{2}$	60	3000	4.0	0.03
DY	$\frac{1}{2}$	60	4000	0.8	0.01
DY	$\frac{1}{2}$	80	200	0.288	0.001
DY	$\frac{1}{2}$	80	500	1.44	0.009
DY	$\frac{1}{2}$	80	1000	2.48	0.02
DY	$\frac{1}{2}$	80	1500	2.29	0.02
DY	$\frac{1}{2}$	80	2000	1.81	0.02

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Mechanism	Spin	Charge [g_D/e]	Mass [GeV]	Efficiency [%]	$\pm$ Error [%]
DY	$\frac{1}{2}$	80	2500	1.29	0.01
DY	$\frac{1}{2}$	80	3000	0.88	0.01
DY	$\frac{1}{2}$	80	4000	0.048	0.002
DY	$\frac{1}{2}$	100	200	0.0513	0.0005
DY	$\frac{1}{2}$	100	500	0.309	0.003
DY	$\frac{1}{2}$	100	1000	0.577	0.007
DY	$\frac{1}{2}$	100	1500	0.505	0.008
DY	$\frac{1}{2}$	100	2000	0.352	0.006
DY	$\frac{1}{2}$	100	2500	0.247	0.005
DY	$\frac{1}{2}$	100	3000	0.119	0.004
DY	$\frac{1}{2}$	100	4000	0.0006	0.0002
PF	0	1	200	4.63	0.0315
PF	0	1	500	14.18	0.0794
PF	0	1	1000	21.99	0.1205
PF	0	1	1500	22.58	0.1309
PF	0	1	2000	19.96	0.1386
PF	0	1	2500	15.81	0.1247
PF	0	1	3000	11.92	0.1078
PF	0	1	4000	6.01	0.0703
PF	0	2	200	0.02	0.0001
PF	0	2	500	0.3	0.0017
PF	0	2	1000	1.22	0.0076

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Mechanism	Spin	Charge [g_D/e]	Mass [GeV]	Efficiency [%]	$\pm$ Error [%]
PF	0	2	1500	1.98	0.0129
PF	0	2	2000	2.34	0.0186
PF	0	2	2500	2.33	0.0196
PF	0	2	3000	1.99	0.018
PF	0	2	4000	0.67	0.0084
PF	0	20	200	4.3	0.0693
PF	0	20	500	6.83	0.0897
PF	0	20	1000	7.26	0.0975
PF	0	20	1500	6.36	0.0922
PF	0	20	2000	5.5	0.0852
PF	0	20	2500	4.45	0.0756
PF	0	20	3000	4.09	0.0705
PF	0	20	4000	2.13	0.0427
PF	0	40	200	7.84	0.0599
PF	0	40	500	14.15	0.0997
PF	0	40	1000	16.02	0.1067
PF	0	40	1500	14.3	0.0976
PF	0	40	2000	11.74	0.0933
PF	0	40	2500	9.92	0.0866
PF	0	40	3000	5.51	0.059
PF	0	40	4000	1.48	0.027
PF	0	60	200	3.35	0.027

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Mechanism	Spin	Charge [g_D/e]	Mass [GeV]	Efficiency [%]	$\pm$ Error [%]
PF	0	60	500	7.05	0.0466
PF	0	60	1000	8.27	0.059
PF	0	60	1500	7.23	0.0561
PF	0	60	2000	5.63	0.0478
PF	0	60	2500	3.77	0.0361
PF	0	60	3000	1.9	0.0256
PF	0	60	4000	0.18	0.0058
PF	0	80	200	0.72	0.0042
PF	0	80	500	2.43	0.018
PF	0	80	1000	3.15	0.0232
PF	0	80	1500	2.53	0.0225
PF	0	80	2000	1.59	0.0176
PF	0	80	2500	0.76	0.0113
PF	0	80	3000	0.28	0.0059
PF	0	80	4000	0.01	0.0006
PF	0	100	200	0.14	0.001
PF	0	100	500	0.65	0.0049
PF	0	100	1000	0.89	0.0076
PF	0	100	1500	0.65	0.0074
PF	0	100	2000	0.32	0.0045
PF	0	100	2500	0.12	0.0025
PF	0	100	3000	0.03	0.001

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Mechanism	Spin	Charge [g_D/e]	Mass [GeV]	Efficiency [%]	$\pm$ Error [%]
PF	0	100	4000	0	0
PF	$\frac{1}{2}$	1	200	3.54	0.0266
PF	$\frac{1}{2}$	1	500	17.16	0.0985
PF	$\frac{1}{2}$	1	1000	30.6	0.1541
PF	$\frac{1}{2}$	1	1500	31.93	0.1642
PF	$\frac{1}{2}$	1	2000	29.1	0.1714
PF	$\frac{1}{2}$	1	2500	23.43	0.1572
PF	$\frac{1}{2}$	1	3000	17.2	0.1357
PF	$\frac{1}{2}$	1	4000	8.39	0.0889
PF	$\frac{1}{2}$	2	200	0.01	0.0001
PF	$\frac{1}{2}$	2	500	0.19	0.0011
PF	$\frac{1}{2}$	2	1000	1.29	0.0083
PF	$\frac{1}{2}$	2	1500	2.91	0.0191
PF	$\frac{1}{2}$	2	2000	3.97	0.0309
PF	$\frac{1}{2}$	2	2500	4.29	0.034
PF	$\frac{1}{2}$	2	3000	3.66	0.0309
PF	$\frac{1}{2}$	2	4000	1.22	0.0142
PF	$\frac{1}{2}$	40	4000	2.49	0.0396
PF	$\frac{1}{2}$	60	200	2.5	0.0225
PF	$\frac{1}{2}$	60	500	7.72	0.0543
PF	$\frac{1}{2}$	60	1000	11.73	0.0825
PF	$\frac{1}{2}$	60	1500	11.59	0.0842

Continued on next page

Mechanism	Spin	Charge [g_D/e]	Mass [GeV]	Efficiency [%]	$\pm$ Error [%]
PF	$\frac{1}{2}$	60	2000	9.74	0.0751
PF	$\frac{1}{2}$	60	2500	6.6	0.058
PF	$\frac{1}{2}$	60	3000	3.52	0.042
PF	$\frac{1}{2}$	60	4000	0.34	0.0097
PF	$\frac{1}{2}$	80	200	0.42	0.0025
PF	$\frac{1}{2}$	80	500	2.18	0.0174
PF	$\frac{1}{2}$	80	1000	3.95	0.0304
PF	$\frac{1}{2}$	80	1500	3.97	0.035
PF	$\frac{1}{2}$	80	2000	2.69	0.0291
PF	$\frac{1}{2}$	80	2500	1.56	0.0208
PF	$\frac{1}{2}$	80	3000	0.62	0.0114
PF	$\frac{1}{2}$	80	4000	0.02	0.0012
PF	$\frac{1}{2}$	100	200	0.08	0.0005
PF	$\frac{1}{2}$	100	500	0.49	0.0038
PF	$\frac{1}{2}$	100	1000	0.98	0.0088
PF	$\frac{1}{2}$	100	1500	0.9	0.0107
PF	$\frac{1}{2}$	100	2000	0.53	0.0075
PF	$\frac{1}{2}$	100	2500	0.21	0.0045
PF	$\frac{1}{2}$	100	3000	0.05	0.0019
PF	$\frac{1}{2}$	100	4000	0	0

Table A.1: Signal efficiencies for spin-0 and spin- $\frac{1}{2}$ HIPs via Drell-Yan (DY) and photon-fusion (PF) [74].

Appendix B

Complete Set of Systematic Uncertainties

Drell-Yan Spin-0 $1g_D$	eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup Up	Pileup Down	Total Up	Total Down
200 GeV	8	-3 $\pm$ 1	12 $\pm$ 0	5 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	13	11
500 GeV	33	-3 $\pm$ 1	6 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	7	5
1000 GeV	52	-3 $\pm$ 1	5 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	6	3
1500 GeV	54	-3 $\pm$ 1	5 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	6	3
2000 GeV	49	-3 $\pm$ 1	5 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	-4 $\pm$ 0	1 $\pm$ 0	-2 $\pm$ 0	2 $\pm$ 0	5	0
2500 GeV	43	-3 $\pm$ 1	5 $\pm$ 0	1 $\pm$ 0	2 $\pm$ 0	-5 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	2 $\pm$ 0	4	-3
3000 GeV	38	-3 $\pm$ 1	5 $\pm$ 0	4 $\pm$ 0	2 $\pm$ 0	-8 $\pm$ 0	2 $\pm$ 0	-3 $\pm$ 0	2 $\pm$ 0	1	-8
4000 GeV	23	-3 $\pm$ 1	10 $\pm$ 0	0 $\pm$ 0	4 $\pm$ 0	-12 $\pm$ 0	3 $\pm$ 0	-4 $\pm$ 0	5 $\pm$ 0	-2	-8

Table B.1: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for Drell-Yan spin-0 monopoles of charge $|g| = 1g_D$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Drell-Yan Spin-0 $2g_D$	eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup		Pileup		Total Up	Total Down
								Up	Down	Up	Down		
200 GeV	0	-16 $\pm$ 6	5 $\pm$ 0	1 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	16	-15
500 GeV	1	-16 $\pm$ 6	6 $\pm$ 0	-4 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	17	-15
1000 GeV	3	-16 $\pm$ 6	5 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	16	-15
1500 GeV	7	-16 $\pm$ 6	5 $\pm$ 0	-1 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	16	-15
2000 GeV	10	-16 $\pm$ 6	5 $\pm$ 0	1 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	16	-15
2500 GeV	12	-16 $\pm$ 6	5 $\pm$ 0	4 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	17	-15
3000 GeV	14	-16 $\pm$ 6	5 $\pm$ 0	3 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	1 $\pm$ 0	1 $\pm$ 0	17	-15
4000 GeV	9	-16 $\pm$ 6	11 $\pm$ 0	2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	2 $\pm$ 0	-2 $\pm$ 0	1 $\pm$ 0	1 $\pm$ 0	1 $\pm$ 0	19	-11

Table B.2: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for Drell-Yan spin-0 monopoles of charge $|g| = 2g_D$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Drell-Yan Spin-0		eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup		Total		
$ z = 20$										Up	Down	Up	Down
200 GeV	6	-5 ± 2	11 ± 0	5 ± 0	7 ± 0	0 ± 0	2 ± 0	-4 ± 0	4 ± 0	15	11		
500 GeV	13	-5 ± 2	15 ± 0	0 ± 0	6 ± 0	0 ± 0	1 ± 0	-4 ± 0	4 ± 0	17	15		
1000 GeV	17	-5 ± 2	7 ± 0	-1 ± 0	8 ± 0	0 ± 0	7 ± 0	-5 ± 0	4 ± 0	14	11		
1500 GeV	15	-5 ± 2	25 ± 0	1 ± 0	7 ± 0	0 ± 0	5 ± 0	-5 ± 0	5 ± 0	28	26		
2000 GeV	14	-5 ± 2	24 ± 0	-1 ± 0	8 ± 0	-2 ± 0	5 ± 0	-6 ± 0	6 ± 0	27	25		
2500 GeV	12	-5 ± 2	11 ± 0	1 ± 0	10 ± 0	-4 ± 0	0 ± 0	-7 ± 0	6 ± 0	16	11		
3000 GeV	11	-5 ± 2	22 ± 0	4 ± 0	11 ± 0	-6 ± 0	9 ± 0	-7 ± 0	7 ± 0	27	24		
4000 GeV	8	-5 ± 2	5 ± 0	0 ± 0	11 ± 0	-17 ± 0	2 ± 0	-8 ± 0	8 ± 0	-6	-15		

Table B.3: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for Drell-Yan spin-0 HECOs of charge $|z| = 20$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Drell-Yan Spin-0 $ z = 40$	eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap		Pileup		Total	
							Up	Down	Up	Down	Up	Down
200 GeV	14	-3 ± 1	5 ± 0	5 ± 0	0 ± 0	0 ± 0	0 ± 0	1 ± 0	-1 ± 0	1 ± 0	7	-2
500 GeV	32	-3 ± 1	5 ± 0	0 ± 0	0 ± 0	0 ± 0	-1 ± 0	1 ± 0	-1 ± 0	1 ± 0	6	4
1000 GeV	40	-3 ± 1	5 ± 0	-1 ± 0	0 ± 0	0 ± 0	-1 ± 0	1 ± 0	-1 ± 0	1 ± 0	6	4
1500 GeV	38	-3 ± 1	5 ± 0	1 ± 0	1 ± 0	-1 ± 0	1 ± 0	2 ± 0	-2 ± 0	2 ± 0	6	4
2000 GeV	34	-3 ± 1	6 ± 0	-1 ± 0	2 ± 0	-3 ± 0	2 ± 0	2 ± 0	-2 ± 0	2 ± 0	6	4
2500 GeV	28	-3 ± 1	7 ± 0	1 ± 0	3 ± 0	-4 ± 0	-2 ± 0	2 ± 0	-2 ± 0	2 ± 0	6	4
3000 GeV	25	-3 ± 1	6 ± 0	4 ± 0	3 ± 0	-5 ± 0	2 ± 0	3 ± 0	-3 ± 0	3 ± 0	7	-2
4000 GeV	13	-3 ± 1	5 ± 0	0 ± 0	4 ± 0	-4 ± 0	4 ± 0	3 ± 0	-3 ± 0	3 ± 0	8	5

Table B.4: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for Drell-Yan spin-0 HECOs of charge $|z| = 40$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Drell-Yan Spin-0		eff.	Material	Birks Hig	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup Up	Pileup Down	Total Up	Total Down
$ z = 60$												
200 GeV	6	-3 $\pm$ 1	5 $\pm$ 0	5 $\pm$ 0	5 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	7	-3
500 GeV	17	-3 $\pm$ 1	5 $\pm$ 0	5 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	5	4
1000 GeV	24	-3 $\pm$ 1	5 $\pm$ 0	5 $\pm$ 0	-1 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	5	4
1500 GeV	23	-3 $\pm$ 1	5 $\pm$ 0	5 $\pm$ 0	1 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	6	4
2000 GeV	20	-3 $\pm$ 1	5 $\pm$ 0	5 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	6	4
2500 GeV	17	-3 $\pm$ 1	5 $\pm$ 0	5 $\pm$ 0	1 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	6	4
3000 GeV	14	-3 $\pm$ 1	5 $\pm$ 0	5 $\pm$ 0	4 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	7	1
4000 GeV	4	-3 $\pm$ 1	5 $\pm$ 0	5 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	-1 $\pm$ 0	2 $\pm$ 0	6	4

Table B.5: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for Drell-Yan spin-0 HECOs of charge $|z| = 60$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Drell-Yan Spin-0 $ z = 80$	eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup		Pileup		Total Up	Total Down
								Up	Down	Up	Down		
200 GeV	1	-6 $\pm$ 3	10 $\pm$ 0	1 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	12	8
500 GeV	6	-6 $\pm$ 3	5 $\pm$ 0	-4 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	9	-6
1000 GeV	9	-6 $\pm$ 3	5 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	8	-5
1500 GeV	9	-6 $\pm$ 3	5 $\pm$ 0	-1 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	8	-5
2000 GeV	7	-6 $\pm$ 3	5 $\pm$ 0	1 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	8	-5
2500 GeV	6	-6 $\pm$ 3	5 $\pm$ 0	4 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	9	-6
3000 GeV	4	-6 $\pm$ 3	5 $\pm$ 0	3 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	9	-5
4000 GeV	0	-6 $\pm$ 3	12 $\pm$ 1	2 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	-4 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	13	9

Table B.6: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for Drell-Yan spin-0 HECOs of charge $|z| = 80$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Drell-Yan Spin-0 $ z = 100$	eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup		Pileup		Total Up	Total Down
								Up	Down	Up	Down		
200 GeV	0	-14 $\pm$ 8	5 $\pm$ 0	1 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	14	-13
500 GeV	1	-14 $\pm$ 8	5 $\pm$ 0	-4 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	15	-13
1000 GeV	2	-14 $\pm$ 8	5 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	14	-13
1500 GeV	2	-14 $\pm$ 8	5 $\pm$ 0	-1 $\pm$ 0	-3 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	14	-13
2000 GeV	2	-14 $\pm$ 8	5 $\pm$ 0	1 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	14	-13
2500 GeV	1	-14 $\pm$ 8	5 $\pm$ 0	4 $\pm$ 0	-3 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	15	-13
3000 GeV	1	-14 $\pm$ 8	5 $\pm$ 0	3 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	15	-13
4000 GeV	0	-14 $\pm$ 8	20 $\pm$ 7	2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	24	15

Table B.7: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for Drell-Yan spin-0 HECOs of charge $|z| = 100$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Drell-Yan Spin- $\frac{1}{2}$ $1g_D$	eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo		Pileup		Pileup		Total Up	Total Down
						Arrival		Up	Down	Up	Down		
200 GeV	3	-4 $\pm$ 1	14 $\pm$ 0	6 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	16	13
500 GeV	12	-4 $\pm$ 1	9 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	10	8
1000 GeV	24	-4 $\pm$ 1	5 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	6	3
1500 GeV	25	-4 $\pm$ 1	5 $\pm$ 0	2 $\pm$ 0	2 $\pm$ 0	-2 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	1 $\pm$ 0	1 $\pm$ 0	7	3
2000 GeV	22	-4 $\pm$ 1	5 $\pm$ 0	0 $\pm$ 0	5 $\pm$ 0	-3 $\pm$ 0	-2 $\pm$ 0	-2 $\pm$ 0	2 $\pm$ 0	2 $\pm$ 0	2 $\pm$ 0	8	5
2500 GeV	19	-4 $\pm$ 1	5 $\pm$ 0	0 $\pm$ 0	7 $\pm$ 0	-5 $\pm$ 0	-3 $\pm$ 0	-3 $\pm$ 0	3 $\pm$ 0	3 $\pm$ 0	3 $\pm$ 0	8	5
3000 GeV	16	-4 $\pm$ 1	5 $\pm$ 0	2 $\pm$ 0	8 $\pm$ 0	-9 $\pm$ 0	-4 $\pm$ 0	-4 $\pm$ 0	4 $\pm$ 0	4 $\pm$ 0	4 $\pm$ 0	7	-4
4000 GeV	8	-4 $\pm$ 1	10 $\pm$ 0	-2 $\pm$ 0	13 $\pm$ 0	-20 $\pm$ 0	-6 $\pm$ 0	-6 $\pm$ 0	6 $\pm$ 0	6 $\pm$ 0	6 $\pm$ 0	-10	-14

Table B.8: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for Drell-Yan spin- $\frac{1}{2}$ monopoles of charge $|g| = 1g_D$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Drell-Yan Spin- $\frac{1}{2}$		eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Pileup		Total	
$2g_D$								Up	Down	Up	Down
200 GeV	0	-18 ± 5	5 ± 0	1 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	19	-17
500 GeV	0	-18 ± 5	6 ± 0	-4 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	19	-17
1000 GeV	1	-18 ± 5	5 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	19	-17
1500 GeV	2	-18 ± 5	6 ± 0	-1 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	19	-17
2000 GeV	2	-18 ± 5	5 ± 0	1 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	19	-17
2500 GeV	3	-18 ± 5	5 ± 0	4 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	19	-18
3000 GeV	4	-18 ± 5	5 ± 0	3 ± 0	0 ± 0	0 ± 0	0 ± 0	-1 ± 0	0 ± 0	19	-18
4000 GeV	2	-18 ± 5	12 ± 0	2 ± 0	2 ± 0	2 ± 0	0 ± 0	-2 ± 0	2 ± 0	22	-14

Table B.9: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for Drell-Yan spin- $\frac{1}{2}$ monopoles of charge $|g| = 2g_D$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Drell-Yan Spin- $\frac{1}{2}$		eff.	Material	Birks High	δ -ray prod.	TRT		Calo Arrival	Pileup		Total	
$ z = 20$						Occ.	Arrival		Up	Down	Up	Down
200 GeV	2	-9 $\pm$ 4	21 $\pm$ 1	6 $\pm$ 0	17 $\pm$ 1	0 $\pm$ 0	-4 $\pm$ 0	4 $\pm$ 0	30	25		
500 GeV	6	-9 $\pm$ 4	18 $\pm$ 0	-1 $\pm$ 0	15 $\pm$ 0	0 $\pm$ 0	-5 $\pm$ 0	5 $\pm$ 0	26	21		
1000 GeV	7	-9 $\pm$ 4	14 $\pm$ 0	-1 $\pm$ 0	16 $\pm$ 0	0 $\pm$ 0	-6 $\pm$ 0	5 $\pm$ 0	24	19		
1500 GeV	6	-9 $\pm$ 4	31 $\pm$ 0	2 $\pm$ 0	19 $\pm$ 0	-2 $\pm$ 0	-6 $\pm$ 0	6 $\pm$ 0	38	34		
2000 GeV	5	-9 $\pm$ 4	25 $\pm$ 1	0 $\pm$ 0	20 $\pm$ 0	-4 $\pm$ 0	-7 $\pm$ 0	7 $\pm$ 0	34	30		
2500 GeV	5	-9 $\pm$ 4	12 $\pm$ 0	0 $\pm$ 0	22 $\pm$ 0	-8 $\pm$ 0	-7 $\pm$ 0	8 $\pm$ 0	27	21		
3000 GeV	4	-9 $\pm$ 4	28 $\pm$ 0	2 $\pm$ 0	22 $\pm$ 0	-14 $\pm$ 0	-8 $\pm$ 0	9 $\pm$ 0	35	30		
4000 GeV	3	-9 $\pm$ 4	15 $\pm$ 0	-2 $\pm$ 0	24 $\pm$ 1	-27 $\pm$ 0	-9 $\pm$ 0	9 $\pm$ 0	15	-11		

Table B.10: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for Drell-Yan spin- $\frac{1}{2}$ HECOs of charge $|z| = 20$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Drell-Yan Spin- $\frac{1}{2}$		eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Pileup Up	Pileup Down	Total Up	Total Down
$ z = 40$											
200 GeV	5	-4 ± 1	5 ± 0	6 ± 0	4 ± 0	0 ± 0	-1 ± 0	1 ± 0		9	-3
500 GeV	13	-4 ± 1	5 ± 0	-1 ± 0	3 ± 0	0 ± 0	-1 ± 0	1 ± 0		7	4
1000 GeV	17	-4 ± 1	6 ± 0	-1 ± 0	3 ± 0	0 ± 0	-1 ± 0	1 ± 0		8	5
1500 GeV	16	-4 ± 1	5 ± 0	2 ± 0	4 ± 0	-1 ± 0	-2 ± 0	2 ± 0		8	4
2000 GeV	13	-4 ± 1	5 ± 0	0 ± 0	5 ± 0	-2 ± 0	-2 ± 0	2 ± 0		8	5
2500 GeV	9	-4 ± 1	9 ± 0	0 ± 0	6 ± 0	-2 ± 0	-3 ± 0	3 ± 0		11	9
3000 GeV	8	-4 ± 1	8 ± 0	2 ± 0	7 ± 0	-3 ± 0	-3 ± 0	3 ± 0		11	8
4000 GeV	4	-4 ± 1	6 ± 0	-2 ± 0	10 ± 0	-5 ± 0	-5 ± 0	-5 ± 0		13	8

Table B.11: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for Drell-Yan spin- $\frac{1}{2}$ HECOs of charge $|z| = 40$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Drell-Yan Spin- $\frac{1}{2}$		eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Pileup Up	Pileup Down	Total Up	Total Down
$ z = 60$											
200 GeV	2		-3 ± 1	5 ± 0	6 ± 0	1 ± 0	0 ± 0	-1 ± 0	1 ± 0	8	-4
500 GeV	5		-3 ± 1	5 ± 0	-1 ± 0	1 ± 0	0 ± 0	0 ± 0	0 ± 0	6	4
1000 GeV	8		-3 ± 1	5 ± 0	-1 ± 0	0 ± 0	0 ± 0	-1 ± 0	1 ± 0	6	3
1500 GeV	8		-3 ± 1	5 ± 0	2 ± 0	1 ± 0	0 ± 0	-1 ± 0	1 ± 0	6	3
2000 GeV	7		-3 ± 1	5 ± 0	0 ± 0	1 ± 0	0 ± 0	-1 ± 0	1 ± 0	6	4
2500 GeV	5		-3 ± 1	5 ± 0	0 ± 0	1 ± 0	0 ± 0	-1 ± 0	1 ± 0	6	4
3000 GeV	4		-3 ± 1	5 ± 0	2 ± 0	2 ± 0	0 ± 0	-1 ± 0	1 ± 0	7	3
4000 GeV	1		-3 ± 1	5 ± 0	-2 ± 0	3 ± 0	0 ± 0	-2 ± 0	2 ± 0	7	3

Table B.12: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for Drell-Yan spin- $\frac{1}{2}$ HECOs of charge $|z| = 60$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Drell-Yan Spin- $\frac{1}{2}$		eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Pileup Up	Pileup Down	Total Up	Total Down
$ z = 80$											
200 GeV	0	-8 $\pm$ 3	11 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	14	7
500 GeV	1	-8 $\pm$ 3	5 $\pm$ 0	-4 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	10	-8
1000 GeV	2	-8 $\pm$ 3	5 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	10	-7
1500 GeV	2	-8 $\pm$ 3	5 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	10	-7
2000 GeV	2	-8 $\pm$ 3	5 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	10	-7
2500 GeV	1	-8 $\pm$ 3	5 $\pm$ 0	4 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	11	-8
3000 GeV	1	-8 $\pm$ 3	5 $\pm$ 0	3 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	10	-7
4000 GeV	0	-8 $\pm$ 3	15 $\pm$ 1	2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	17	12

Table B.13: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for Drell-Yan spin- $\frac{1}{2}$ HECOs of charge $|z| = 80$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Drell-Yan Spin- $\frac{1}{2}$		eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo		Pileup		Total	
$ z = 100$							Arrival	Up	Down	Up	Down	
200 GeV	0	-16 ± 13	5 ± 0	1 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	17	-15
500 GeV	0	-16 ± 13	5 ± 0	-4 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	17	-15
1000 GeV	1	-16 ± 13	5 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	17	-15
1500 GeV	1	-16 ± 13	5 ± 0	-1 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	17	-15
2000 GeV	0	-16 ± 13	5 ± 0	1 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	17	-15
2500 GeV	0	-16 ± 13	5 ± 0	4 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	-1 ± 0	1 ± 0	17	-16
3000 GeV	0	-16 ± 13	5 ± 0	3 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	-1 ± 0	1 ± 0	17	-15
4000 GeV	0	-16 ± 13	5 ± 0	2 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	-3 ± 0	4 ± 0	17	-15

Table B.14: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for Drell-Yan spin- $\frac{1}{2}$ HECOs of charge $|z| = 100$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Photon Fusion Spin-0 $1g_D$	eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup		Total	
								Up	Down	Up	Down
200 GeV	5	-3 $\pm$ 1	11 $\pm$ 0	5 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	12	9
500 GeV	14	-3 $\pm$ 1	7 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	7	6
1000 GeV	22	-3 $\pm$ 1	5 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	6	3
1500 GeV	23	-3 $\pm$ 1	5 $\pm$ 0	2 $\pm$ 0	0 $\pm$ 0	-4 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	6	-2
2000 GeV	20	-3 $\pm$ 1	6 $\pm$ 0	1 $\pm$ 0	2 $\pm$ 0	-8 $\pm$ 0	1 $\pm$ 0	-2 $\pm$ 0	2 $\pm$ 0	-4	-7
2500 GeV	16	-3 $\pm$ 1	6 $\pm$ 0	0 $\pm$ 0	4 $\pm$ 0	-10 $\pm$ 0	0 $\pm$ 0	-3 $\pm$ 0	3 $\pm$ 0	-6	-9
3000 GeV	12	-3 $\pm$ 1	6 $\pm$ 0	3 $\pm$ 0	5 $\pm$ 0	-16 $\pm$ 0	2 $\pm$ 0	-4 $\pm$ 0	4 $\pm$ 0	-12	-15
4000 GeV	6	-3 $\pm$ 1	14 $\pm$ 0	-7 $\pm$ 0	5 $\pm$ 0	-30 $\pm$ 0	3 $\pm$ 0	-6 $\pm$ 0	6 $\pm$ 0	-24	-28

Table B.15: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for photon fusion spin-0 monopoles of charge $|g| = 1g_D$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Photon Fusion Spin-0		eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup		Total	
$2g_D$									Up	Down	Up	Down
200 GeV	0	-15 ± 6	5 ± 0	1 ± 0	-2 ± 0	0 ± 0	2 ± 0	0 ± 0	0 ± 0	16	-14	
500 GeV	0	-15 ± 6	6 ± 0	-4 ± 0	-2 ± 0	0 ± 0	2 ± 0	0 ± 0	0 ± 0	16	-14	
1000 GeV	1	-15 ± 6	5 ± 0	0 ± 0	-2 ± 0	0 ± 0	1 ± 0	0 ± 0	0 ± 0	16	-14	
1500 GeV	2	-15 ± 6	5 ± 0	-1 ± 0	-2 ± 0	0 ± 0	2 ± 0	0 ± 0	0 ± 0	16	-14	
2000 GeV	2	-15 ± 6	5 ± 0	1 ± 0	-2 ± 0	0 ± 0	1 ± 0	0 ± 0	0 ± 0	16	-14	
2500 GeV	2	-15 ± 6	5 ± 0	4 ± 0	-1 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	16	-15	
3000 GeV	2	-15 ± 6	5 ± 0	5 ± 0	-2 ± 0	0 ± 0	0 ± 0	-1 ± 0	1 ± 0	16	-15	
4000 GeV	1	-15 ± 6	16 ± 0	2 ± 0	0 ± 0	0 ± 0	2 ± 0	-2 ± 0	2 ± 0	22	5	

Table B.16: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for photon fusion spin-0 monopoles of charge $|g| = 2g_D$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Photon Fusion Spin-0 $ z = 20$	eff.	Material	Birk High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup		Pileup		Total	
								Up	Down	Up	Down	Up	Down
200 GeV	4	-10 $\pm$ 5	10 $\pm$ 0	5 $\pm$ 0	7 $\pm$ 0	0 $\pm$ 0	2 $\pm$ 0	-4 $\pm$ 0	4 $\pm$ 0	17	4		
500 GeV	7	-10 $\pm$ 5	16 $\pm$ 0	0 $\pm$ 0	6 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	-4 $\pm$ 0	4 $\pm$ 0	20	14		
1000 GeV	7	-10 $\pm$ 5	12 $\pm$ 0	-1 $\pm$ 0	9 $\pm$ 0	-1 $\pm$ 0	7 $\pm$ 0	-5 $\pm$ 0	5 $\pm$ 0	20	12		
1500 GeV	6	-10 $\pm$ 5	29 $\pm$ 1	2 $\pm$ 0	7 $\pm$ 0	-2 $\pm$ 0	5 $\pm$ 0	-6 $\pm$ 0	6 $\pm$ 0	32	28		
2000 GeV	6	-10 $\pm$ 5	25 $\pm$ 1	1 $\pm$ 0	8 $\pm$ 0	-7 $\pm$ 0	5 $\pm$ 0	-7 $\pm$ 0	7 $\pm$ 0	29	23		
2500 GeV	4	-10 $\pm$ 5	12 $\pm$ 0	0 $\pm$ 0	11 $\pm$ 0	-12 $\pm$ 0	0 $\pm$ 0	-8 $\pm$ 0	7 $\pm$ 0	16	-6		
3000 GeV	4	-10 $\pm$ 5	29 $\pm$ 1	3 $\pm$ 0	11 $\pm$ 0	-19 $\pm$ 0	9 $\pm$ 0	-8 $\pm$ 0	9 $\pm$ 0	29	23		
4000 GeV	2	-10 $\pm$ 5	39 $\pm$ 2	-7 $\pm$ 0	14 $\pm$ 0	-40 $\pm$ 1	2 $\pm$ 0	-9 $\pm$ 0	9 $\pm$ 0	20	-5		

Table B.17: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for photon fusion spin-0 HECOs of charge $|z| = 20$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Photon Fusion Spin-0		eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup		Pileup		Total Up	Total Down
$ z = 40$									Up	Down				
200 GeV	8	-4 $\pm$ 1	5 $\pm$ 0	5 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	1 $\pm$ 0	8	-4	
500 GeV	14	-4 $\pm$ 1	5 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	1 $\pm$ 0	7	2	
1000 GeV	16	-4 $\pm$ 1	5 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	1 $\pm$ 0	7	2	
1500 GeV	14	-4 $\pm$ 1	5 $\pm$ 0	2 $\pm$ 0	1 $\pm$ 0	1 $\pm$ 0	-3 $\pm$ 0	1 $\pm$ 0	-2 $\pm$ 0	2 $\pm$ 0	2 $\pm$ 0	7	-3	
2000 GeV	12	-4 $\pm$ 1	5 $\pm$ 0	1 $\pm$ 0	2 $\pm$ 0	2 $\pm$ 0	-7 $\pm$ 0	2 $\pm$ 0	-2 $\pm$ 0	2 $\pm$ 0	2 $\pm$ 0	4	-6	
2500 GeV	10	-4 $\pm$ 1	10 $\pm$ 0	0 $\pm$ 0	4 $\pm$ 0	4 $\pm$ 0	-9 $\pm$ 0	-2 $\pm$ 0	-2 $\pm$ 0	2 $\pm$ 0	2 $\pm$ 0	8	3	
3000 GeV	6	-4 $\pm$ 1	8 $\pm$ 0	3 $\pm$ 0	4 $\pm$ 0	4 $\pm$ 0	-10 $\pm$ 0	2 $\pm$ 0	-3 $\pm$ 0	3 $\pm$ 0	3 $\pm$ 0	5	-7	
4000 GeV	1	-4 $\pm$ 1	9 $\pm$ 0	-7 $\pm$ 0	5 $\pm$ 0	5 $\pm$ 0	-7 $\pm$ 0	4 $\pm$ 0	-3 $\pm$ 0	4 $\pm$ 0	4 $\pm$ 0	12	-4	

Table B.18: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for photon fusion spin-0 HECOs of charge $|z| = 40$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Photon Fusion Spin-0 $ z = 60$	eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap		Pileup		Total	
							Up	Down	Up	Down	Up	Down
200 GeV	3	-3 ± 1	5 ± 0	5 ± 0	-2 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	7	-3
500 GeV	7	-3 ± 1	5 ± 0	0 ± 0	-2 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	1 ± 0	6	3
1000 GeV	8	-3 ± 1	5 ± 0	-1 ± 0	-2 ± 0	0 ± 0	-1 ± 0	-1 ± 0	-1 ± 0	1 ± 0	6	3
1500 GeV	7	-3 ± 1	5 ± 0	2 ± 0	-2 ± 0	0 ± 0	-1 ± 0	-1 ± 0	-1 ± 0	1 ± 0	6	2
2000 GeV	6	-3 ± 1	5 ± 0	1 ± 0	-1 ± 0	0 ± 0	1 ± 0	-1 ± 0	-1 ± 0	1 ± 0	6	3
2500 GeV	4	-3 ± 1	5 ± 0	0 ± 0	-1 ± 0	-1 ± 0	1 ± 0	-1 ± 0	-1 ± 0	1 ± 0	6	3
3000 GeV	2	-3 ± 1	5 ± 0	3 ± 0	0 ± 0	-2 ± 0	0 ± 0	-1 ± 0	-1 ± 0	1 ± 0	7	-2
4000 GeV	0	-3 ± 1	5 ± 0	-7 ± 0	3 ± 0	-2 ± 0	1 ± 0	-1 ± 0	-1 ± 0	2 ± 0	10	-6

Table B.19: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for photon fusion spin-0 HECOs of charge $|z| = 60$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Photon Fusion Spin-0		eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup		Total	
$ z = 80$									Up	Down	Up	Down
200 GeV	1	-7 ± 4	10 ± 0	1 ± 0	-2 ± 0	0 ± 0	1 ± 0	0 ± 0	0 ± 0	12	6	
500 GeV	2	-7 ± 4	5 ± 0	-4 ± 0	-2 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	9	-7	
1000 GeV	3	-7 ± 4	5 ± 0	0 ± 0	-2 ± 0	0 ± 0	1 ± 0	0 ± 0	0 ± 0	9	-6	
1500 GeV	3	-7 ± 4	5 ± 0	-1 ± 0	-2 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	9	-6	
2000 GeV	2	-7 ± 4	5 ± 0	1 ± 0	-2 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	9	-6	
2500 GeV	1	-7 ± 4	5 ± 0	4 ± 0	-2 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	10	-7	
3000 GeV	0	-7 ± 4	5 ± 0	5 ± 0	-3 ± 0	0 ± 0	1 ± 0	0 ± 0	1 ± 0	10	-7	
4000 GeV	0	-7 ± 4	17 ± 2	2 ± 0	-5 ± 1	0 ± 0	-4 ± 0	-5 ± 1	0 ± 0	18	14	

Table B.20: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for photon fusion spin-0 HECOs of charge $|z| = 80$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Photon Fusion Spin-0 $ z = 100$	eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap		Pileup		Pileup		Total	Total
													Up	Down
200 GeV	0	-14 ± 9	5 ± 0	1 ± 0	-2 ± 0	0 ± 0	-1 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	14	-13
500 GeV	1	-14 ± 9	5 ± 0	-4 ± 0	-2 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	15	-13
1000 GeV	1	-14 ± 9	5 ± 0	0 ± 0	-2 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	14	-13
1500 GeV	1	-14 ± 9	5 ± 0	-1 ± 0	-3 ± 0	0 ± 0	1 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	14	-13
2000 GeV	0	-14 ± 9	5 ± 0	1 ± 0	-2 ± 0	0 ± 0	1 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	14	-13
2500 GeV	0	-14 ± 9	7 ± 0	4 ± 0	-2 ± 0	0 ± 0	1 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	16	-13
3000 GeV	0	-14 ± 9	5 ± 0	5 ± 0	-1 ± 0	-1 ± 0	0 ± 0	-1 ± 0	-1 ± 0	-1 ± 0	-1 ± 0	-1 ± 0	15	-14
4000 GeV	0	-14 ± 9	5 ± 0	2 ± 0	0 ± 0	0 ± 0	1 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	15	-13

Table B.21: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for photon fusion spin-0 HECOs of charge $|z| = 100$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Photon Fusion Spin- $\frac{1}{2}$		eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup		Total	
$1g_D$									Up	Down	Up	Down
200 GeV	4	-4 $\pm$ 1	14 $\pm$ 0	6 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	15	12	
500 GeV	17	-4 $\pm$ 1	8 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	8	7	
1000 GeV	31	-4 $\pm$ 1	5 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	6	3	
1500 GeV	32	-4 $\pm$ 1	5 $\pm$ 0	2 $\pm$ 0	0 $\pm$ 0	-3 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	6	-1	
2000 GeV	29	-4 $\pm$ 1	5 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	-7 $\pm$ 0	1 $\pm$ 0	-2 $\pm$ 0	2 $\pm$ 0	-2	-6	
2500 GeV	23	-4 $\pm$ 1	5 $\pm$ 0	0 $\pm$ 0	3 $\pm$ 0	-9 $\pm$ 0	0 $\pm$ 0	-3 $\pm$ 0	2 $\pm$ 0	-5	-8	
3000 GeV	17	-4 $\pm$ 1	5 $\pm$ 0	4 $\pm$ 0	4 $\pm$ 0	-13 $\pm$ 0	2 $\pm$ 0	-4 $\pm$ 0	3 $\pm$ 0	-10	-13	
4000 GeV	8	-4 $\pm$ 1	14 $\pm$ 0	-6 $\pm$ 0	5 $\pm$ 0	-27 $\pm$ 0	3 $\pm$ 0	-6 $\pm$ 0	6 $\pm$ 0	-21	-25	

Table B.22: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for photon fusion spin- $\frac{1}{2}$ monopoles of charge $|g| = 1g_D$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Photon Fusion Spin- $\frac{1}{2}$		eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup		Total	
$2g_D$									Up	Down	Up	Down
200 GeV	0	-17 $\pm$ 6	5 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	18	-16
500 GeV	0	-17 $\pm$ 6	6 $\pm$ 0	-4 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	18	-17
1000 GeV	1	-17 $\pm$ 6	5 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	18	-16
1500 GeV	3	-17 $\pm$ 6	5 $\pm$ 0	-1 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	18	-16
2000 GeV	4	-17 $\pm$ 6	5 $\pm$ 0	1 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	18	-17
2500 GeV	4	-17 $\pm$ 6	5 $\pm$ 0	4 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	18	-17
3000 GeV	4	-17 $\pm$ 6	5 $\pm$ 0	4 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	18	-17
4000 GeV	1	-17 $\pm$ 6	14 $\pm$ 0	3 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	2 $\pm$ 0	-2 $\pm$ 0	2 $\pm$ 0	2 $\pm$ 0	23	-10

Table B.23: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for photon fusion spin- $\frac{1}{2}$ monopoles of charge $|g| = 2g_D$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Photon Fusion Spin- $\frac{1}{2}$ $ z = 20$	eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup		Total	
								Up	Down	Up	Down
200 GeV	4	-8 $\pm$ 4	12 $\pm$ 0	6 $\pm$ 0	8 $\pm$ 0	0 $\pm$ 0	2 $\pm$ 0	-4 $\pm$ 0	4 $\pm$ 0	18	9
500 GeV	8	-8 $\pm$ 4	16 $\pm$ 0	0 $\pm$ 0	6 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	-4 $\pm$ 0	4 $\pm$ 0	20	15
1000 GeV	9	-8 $\pm$ 4	14 $\pm$ 0	-1 $\pm$ 0	8 $\pm$ 0	0 $\pm$ 0	7 $\pm$ 0	-5 $\pm$ 0	5 $\pm$ 0	20	15
1500 GeV	9	-8 $\pm$ 4	26 $\pm$ 0	2 $\pm$ 0	8 $\pm$ 0	-2 $\pm$ 0	5 $\pm$ 0	-5 $\pm$ 0	5 $\pm$ 0	30	26
2000 GeV	8	-8 $\pm$ 4	22 $\pm$ 0	0 $\pm$ 0	9 $\pm$ 0	-5 $\pm$ 0	5 $\pm$ 0	-7 $\pm$ 0	6 $\pm$ 0	25	21
2500 GeV	6	-8 $\pm$ 4	8 $\pm$ 0	0 $\pm$ 0	10 $\pm$ 0	-8 $\pm$ 0	0 $\pm$ 0	-7 $\pm$ 0	7 $\pm$ 0	15	-5
3000 GeV	6	-8 $\pm$ 4	26 $\pm$ 1	4 $\pm$ 0	11 $\pm$ 0	-14 $\pm$ 0	9 $\pm$ 0	-7 $\pm$ 0	9 $\pm$ 0	29	24
4000 GeV	3	-8 $\pm$ 4	26 $\pm$ 1	-6 $\pm$ 0	14 $\pm$ 0	-34 $\pm$ 1	2 $\pm$ 0	-8 $\pm$ 0	9 $\pm$ 0	-9	-21

Table B.24: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for photon fusion spin- $\frac{1}{2}$ HECOs of charge $|z| = 20$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Photon Fusion Spin- $\frac{1}{2}$		eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup		Total	
$ z = 40$									Up	Down	Up	Down
200 GeV	7	-4 ± 1	5 ± 0	6 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	-1 ± 0	1 ± 0	9	-5
500 GeV	17	-4 ± 1	5 ± 0	0 ± 0	0 ± 0	0 ± 0	0 ± 0	-1 ± 0	-1 ± 0	1 ± 0	6	3
1000 GeV	22	-4 ± 1	5 ± 0	-1 ± 0	0 ± 0	0 ± 0	0 ± 0	-1 ± 0	-1 ± 0	1 ± 0	6	3
1500 GeV	22	-4 ± 1	5 ± 0	2 ± 0	1 ± 0	-2 ± 0	1 ± 0	1 ± 0	-2 ± 0	2 ± 0	7	1
2000 GeV	18	-4 ± 1	5 ± 0	0 ± 0	2 ± 0	-5 ± 0	2 ± 0	2 ± 0	-2 ± 0	2 ± 0	5	-4
2500 GeV	15	-4 ± 1	8 ± 0	0 ± 0	4 ± 0	-9 ± 0	-2 ± 0	-2 ± 0	-2 ± 0	2 ± 0	5	-4
3000 GeV	9	-4 ± 1	9 ± 0	4 ± 0	4 ± 0	-8 ± 0	2 ± 0	2 ± 0	-3 ± 0	3 ± 0	8	-4
4000 GeV	2	-4 ± 1	9 ± 0	-6 ± 0	4 ± 0	-7 ± 0	4 ± 0	4 ± 0	-4 ± 0	4 ± 0	11	1

Table B.25: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for photon fusion spin- $\frac{1}{2}$ HECOs of charge $|z| = 40$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Photon Fusion Spin- $\frac{1}{2}$ $ z = 60$	eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup Up	Pileup Down	Total Up	Total Down
200 GeV	3	-4 $\pm$ 1	5 $\pm$ 0	6 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	8	-5
500 GeV	8	-4 $\pm$ 1	5 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	6	3
1000 GeV	12	-4 $\pm$ 1	5 $\pm$ 0	-1 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	6	3
1500 GeV	12	-4 $\pm$ 1	5 $\pm$ 0	2 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	6	2
2000 GeV	10	-4 $\pm$ 1	5 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	6	3
2500 GeV	7	-4 $\pm$ 1	5 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	6	3
3000 GeV	4	-4 $\pm$ 1	5 $\pm$ 0	4 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	7	-2
4000 GeV	0	-4 $\pm$ 1	5 $\pm$ 0	-6 $\pm$ 0	1 $\pm$ 0	-1 $\pm$ 0	1 $\pm$ 0	-2 $\pm$ 0	1 $\pm$ 0	9	-4

Table B.26: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for photon fusion spin- $\frac{1}{2}$ HECOs of charge $|z| = 60$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Photon Fusion Spin- $\frac{1}{2}$		eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup		Pileup		Total Up	Total Down
$ z = 80$									Up	Down				
200 GeV		0	-8 $\pm$ 4	10 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	13	6	
500 GeV		2	-8 $\pm$ 4	5 $\pm$ 0	-4 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	10	-7	
1000 GeV		4	-8 $\pm$ 4	5 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	9	-6	
1500 GeV		4	-8 $\pm$ 4	5 $\pm$ 0	-1 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	9	-6	
2000 GeV		3	-8 $\pm$ 4	5 $\pm$ 0	1 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	9	-6	
2500 GeV		2	-8 $\pm$ 4	5 $\pm$ 0	4 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	10	-7	
3000 GeV		1	-8 $\pm$ 4	5 $\pm$ 0	4 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	10	-7	
4000 GeV		0	-8 $\pm$ 4	11 $\pm$ 1	3 $\pm$ 0	-3 $\pm$ 0	0 $\pm$ 0	-4 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	13	6	

Table B.27: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for photon fusion spin- $\frac{1}{2}$ HECOs of charge $|z| = 80$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

Photon Fusion Spin- $\frac{1}{2}$		eff.	Material	Birks High	δ -ray prod.	TRT Occ.	Calo Arrival	Extrap	Pileup Up	Pileup Down	Total Up	Total Down
$ z = 100$												
200 GeV	0	-14 $\pm$ 9	5 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	15	-13
500 GeV	0	-14 $\pm$ 9	5 $\pm$ 0	-4 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	15	-14
1000 GeV	1	-14 $\pm$ 9	5 $\pm$ 0	0 $\pm$ 0	-2 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	15	-13
1500 GeV	1	-14 $\pm$ 9	5 $\pm$ 0	-1 $\pm$ 0	-3 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	15	-13
2000 GeV	1	-14 $\pm$ 9	5 $\pm$ 0	1 $\pm$ 0	-3 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	15	-13
2500 GeV	0	-14 $\pm$ 9	8 $\pm$ 0	4 $\pm$ 0	-3 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	-1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	16	-13
3000 GeV	0	-14 $\pm$ 9	5 $\pm$ 0	4 $\pm$ 0	-2 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	15	-14
4000 GeV	0	-14 $\pm$ 9	5 $\pm$ 0	3 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	1 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	0 $\pm$ 0	15	-13

Table B.28: Relative uncertainties (sys $\pm$ stat) on the signal efficiencies in percentages for photon fusion spin- $\frac{1}{2}$ HECOs of charge $|z| = 100$. The total relative uncertainties are calculated as quadratic sums of the individual relative uncertainties. Note that the detector-material uncertainty and the δ -ray uncertainty are taken as symmetric, as described in Section 6.2.

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