

1 A search for sparticles in zero lepton final states

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## ABSTRACT

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A search for sparticles in zero lepton final states

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Russell W. Smith

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*Acknowledgements*



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*Dedication*



*Introduction*

74 Particle physics is a remarkably successful field of scientific inquiry. The ability to  
 75 precisely predict the properties of a exceedingly wide range of physical phenomena,  
 76 such as the description of the cosmic microwave background [1, 2], the understanding  
 77 of the anomalous magnetic dipole moment of the electron [3, 4], and the measurement  
 78 of the number of weakly-interacting neutrino flavors [5] is truly amazing.

79 The theory that has allowed this range of predictions is the *Standard Model*  
 80 of particle physics (SM). The Standard Model combines the electroweak theory of  
 81 Glashow, Weinberg, and Salam [6–8] with the theory of the strong interactions, as  
 82 first envisioned by Gell-Mann and Zweig [9, 10]. This quantum field theory (QFT)  
 83 contains a tiny number of particles, whose interactions describe phenomena up to at  
 84 least the TeV scale. These particles are manifestations of the fields of the Standard  
 85 Model, after application of the Higgs Mechanism. The particle content of the SM  
 86 consists only of the six quarks, the six leptons, the four gauge bosons, and the scalar  
 87 Higgs boson.

88 Despite its impressive range of described phenomena, the Standard Model has  
 89 some theoretical and experimental deficiencies. The SM contains 26 free parameters  
 90 <sup>1</sup>. It would be more theoretically pleasing to understand these free parameters in  
 91 terms of a more fundamental theory. The major theoretical concern of the Standard  
 92 Model, as it pertains to this thesis, is the *hierarchy problem*[11–15]. The light mass

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<sup>1</sup>This is the Standard Model corrected to include neutrino masses. These parameters are the fermion masses (6 leptons, 6 quarks), CKM and PMNS mixing angles (8 angles, 2 CP-violating phases), W/Z/Higgs masses (3), the Higgs field expectation value, and the couplings of the strong, weak, and electromagnetic forces (3  $\alpha_{force}$  ).

of the Higgs boson (125 GeV) should be quadratically dependent on the scale of UV physics, due to the quantum corrections from high-energy physics processes. The most perplexing experimental issue is the existence of *dark matter*, as demonstrated by galactic rotation curves [16–22]. This data has shown that there exists additional matter which has not yet been seen interacting with the particles of the Standard Model. There is no particle in the SM which can act as a candidate for dark matter.

Both of these major issues, as well as numerous others, can be solved by the introduction of *supersymmetry* (SUSY) [15, 23–35]. In supersymmetric theories, each SM particles has a so-called *superpartner*, or sparticle partner, differing from given SM particle by 1/2 in spin. These theories solve the hierarchy problem, since the quantum corrections induced from the superpartners exactly cancel those induced by the SM particles. In addition, these theories are usually constructed assuming *R*–parity, which can be thought of as the “charge” of supersymmetry, with SM particles having  $R = 1$  and sparticles having  $R = -1$ . In collider experiments, since the incoming SM particles have total  $R = 1$ , the resulting sparticles are produced in pairs. This produces a rich phenomenology, which is characterized by significant hadronic activity and large missing transverse energy ( $E_T^{\text{miss}}$ ), which provide significant discrimination against SM backgrounds [36].

Despite the power of searches for supersymmetry where  $E_T^{\text{miss}}$  is a primary discriminating variable, there has been significant interest in the use of other variables to discriminate against SM backgrounds. These include searches employing variables such as  $\alpha T$ ,  $M_{T,2}$ , and the razor variables ( $M_R, R^2$ ) [37–47]. In this thesis, we will present the first search for supersymmetry using the novel Recursive Jigsaw Reconstruction (RJR) technique. RJR can be considered the conceptual successor of the razor variables. We impose a particular final state “decay tree” on an events, which roughly corresponds to a simplified Feynmann diagram in decays containing weakly-interacting particles. We account for the missing degrees of freedom associated

120 to the weakly-interacting particles by a series of simplifying assumptions, which allow  
121 us to calculate our variables of interest at each step in the decay tree. This allows an  
122 unprecedented understanding of the internal structure of the decay and the ability to  
123 construct additional variables to reject Standard Model backgrounds.

124 This thesis details a search for the superpartners of the gluon and quarks, the  
125 gluino and squarks, in final states with zero leptons, with  $13.3 \text{ fb}^{-1}$  of data using the  
126 ATLAS detector. We organize the thesis as follows. The theoretical foundations of  
127 the Standard Model and supersymmetry are described in Chapters 2 and 3. The  
128 Large Hadron Collider and the ATLAS detector are presented in Chapters 4 and 5.  
129 Chapter 5 provides a detailed description of Recursive Jigsaw Reconstruction and a  
130 description of the variables used for the particular search presented in this thesis.  
131 Chapter 6 presents the details of the analysis, including details of the dataset, object  
132 reconstruction, and selections used. In Chapter 7, the final results are presented;  
133 since there is no evidence of a supersymmetric signal in the analysis, we present the  
134 final exclusion curves in simplified supersymmetric models.



137 **2.1 Overview**

138 A Standard Model is another name for a theory of the internal symmetry group  
 139  $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ , with its associated set of parameters. *The Standard*  
 140 Model refers specifically to a Standard Model with the proper parameters to describe  
 141 the universe. The SM is the culmination of years of work in both theoretical  
 142 and experimental particle physics. In this thesis, we take the view that theorists cite  
 143 construct a model with the field content and symmetries as inputs, and write down the  
 144 most general Lagrangian consistent with those symmetries. Assuming this model is  
 145 compatible with nature (in particular, the predictions of the model are consistent with  
 146 previous experiments), experimentalists are responsible measuring the parameters of  
 147 this model. This will be applicable for this chapter and the following one.

148 Additional theoretical background is in 10.6. The philosophy and notations are  
 149 inspired by [48, 49].

150 **2.2 Field Content**

The Standard Model field content is

$$\begin{aligned} \text{Fermions} &: Q_L(3, 2)_{+1/3}, U_R(3, 1)_{+4/3}, D_R(3, 1)_{-2/3}, L_L(1, 2)_{-1}, E_R(1, 1)_{-2} \\ \text{Scalar (Higgs)} &: \phi(1, 2)_{+1} \\ \text{Vector Fields} &: G^\mu(8, 1)_0, W^\mu(1, 3)_0, B^\mu(1, 1)_0 \end{aligned} \tag{2.1}$$

151 where the  $(A, B)_Y$  notation represents the irreducible representation under  $SU(3)$   
152 and  $SU(2)$ , with  $Y$  being the electroweak hypercharge. Each of these fermion fields  
153 has an additional index, representing the three generation of fermions.

154 We observed that  $Q_L, U_R$ , and  $D_R$  are triplets under  $SU(3)_C$ ; these are the *quark*  
155 fields. The *color* group,  $SU(3)_C$  is mediated by the *gluon* field  $G^\mu(8, 1)_0$ , which has  
156 8 degrees of freedom. The fermion fields  $L_L(1, 2)_{-1}$  and  $E_R(1, 1)_{-2}$  are singlets under  
157  $SU(3)_C$ ; we call them the *lepton* fields.

158 Next, we note the “left-handed” (“right-handed”) fermion fields, denoted by  $L$  ( $R$ )  
159 subscript, The left-handed fields form doublets under  $SU(2)_L$ . These are mediated  
160 by the three degrees of freedom of the “W” fields  $W^\mu(1, 3)_0$ . These fields only act  
161 on the left-handed particles of the Standard Model. This is the reflection of the  
162 “chirality” of the Standard Model; the left-handed and right-handed particles are  
163 treated differently by the electroweak forces. The right-handed fields,  $U_R, D_R$ , and  
164  $E_R$ , are singlets under  $SU(2)_L$ .

165 The  $U(1)_Y$  symmetry is associated to the  $B^\mu(1, 1)_0$  boson with one degree of  
166 freedom. The charge  $Y$  is known as the electroweak hypercharge.

167 To better understand the phenomenology of the Standard Model, let us investigate  
168 each of the *sectors* of the Standard Model separately.

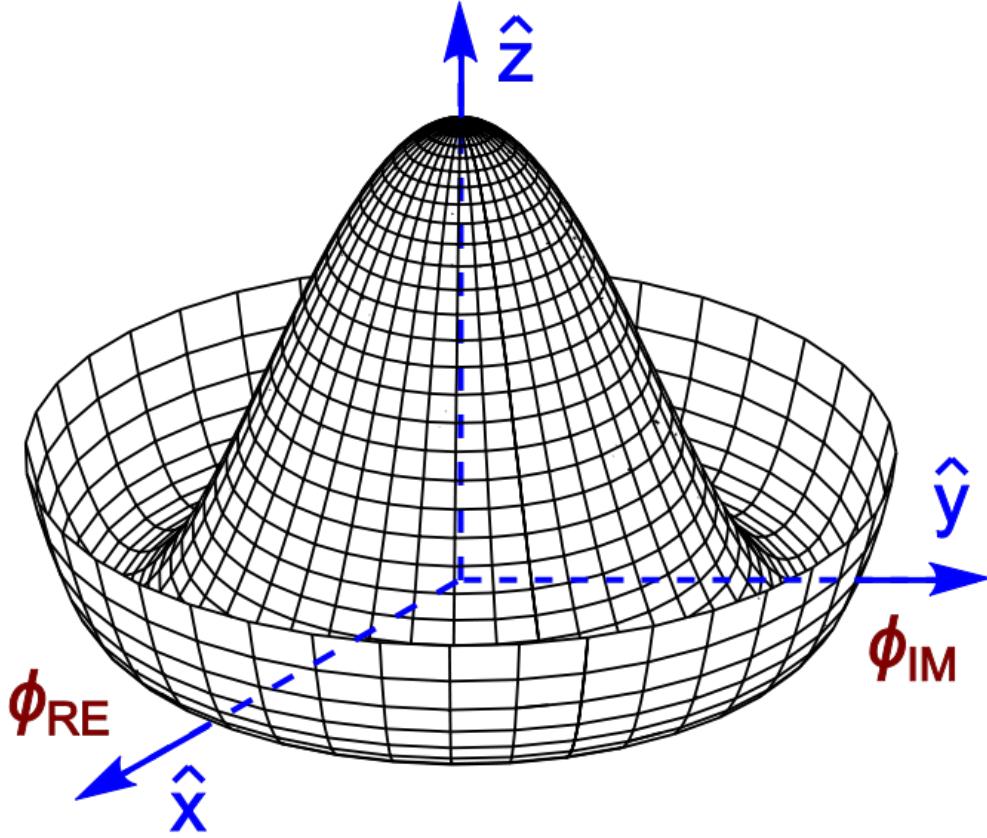
## 169 Electroweak sector

The electroweak sector refers to the  $SU(2)_L \otimes U(1)_Y$  portion of the Standard  
Model gauge group. Following our philosophy of writing all gauge-invariant and  
renormalizable terms, the electroweak Lagrangian can be written as

$$\mathcal{L} = W_a^{\mu\nu} W_{\mu\nu}^a + B^{\mu\nu} B_{\mu\nu} + (D^\mu \phi)^\dagger D_\mu \phi - \mu^2 \phi^\dagger \phi - \lambda (\phi^\dagger \phi)^2. \quad (2.2)$$

where  $W_a^{\mu\nu}$  are the three ( $a = 1, 2, 3$ ) gauge bosons associated to the  $SU(2)_L$  gauge  
group,  $B^{\mu\nu}$  is the one gauge boson of the  $U(1)_Y$  gauge group, and  $\phi$  is the complex

Figure 2.1: Sombrero potential



Higgs multiplet. The covariant derivative  $D^\mu$  is given by

$$D^\mu = \partial^\mu + \frac{ig}{2} W_a^\mu \sigma_a + \frac{ig'}{2} B^\mu \quad (2.3)$$

where  $i\sigma_a$  are the Pauli matrices times the imaginary constant, which are the generators for  $SU(2)_L$ , and  $g$  and  $g'$  are the  $SU(2)_L$  and  $U(1)_Y$  coupling constants, respectively. The field strength tensors  $W_a^{\mu\nu}$  and  $B^{\mu\nu}$  are given by the commutator of the covariant derivative associated to each field

$$B^{\mu\nu} = \partial^\mu B^\nu - \partial^\nu B^\mu \quad (2.4)$$

$$W_a^{\mu\nu} = \partial^\mu W_a^\nu - \partial^\nu W_a^\mu - g\epsilon_{abc} W_a^\mu W_b^\nu, \quad i = 1, 2, 3$$

171      The terms in the Lagrangian 2.2 proportional to  $\mu^2$  and  $\lambda$  make up the “Higgs  
 172 potential” [50]. As normal (see Appendix 10.6), we restrict  $\lambda > 0$  to guarantee our  
 173 potential is bounded from below, and we also require  $\mu^2 < 0$ , which gives us the  
 174 standard “sombrero” potential shown in 2.1.

This potential has infinitely many minima at  $\langle \phi \rangle = \sqrt{2m/\lambda}$ ; the ground state is *spontaneously* broken by the choice of ground state, which induces a vacuum expectation value (VEV). Without loss of generality, we can choose the Higgs field  $\phi$  to point in the real direction, and write the Higgs field  $\phi$  in the following form :

$$\phi = \frac{1}{\sqrt{2}} \exp\left(\frac{i}{v} \sigma_a \theta_a\right) \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix}. \quad (2.5)$$

We choose a gauge to rotate away the dependence on  $\theta_a$ , such that we can write simply

$$\phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix}. \quad (2.6)$$

Now, we can see how the masses of the vector bosons are generated from the application of the Higgs mechanism. We plug Eq.2.6 back into the electroweak Lagrangian, and only showing the relevant mass terms in the vacuum state where  $h(x) = 0$  see that (dropping the Lorentz indices) :

$$\begin{aligned} \mathcal{L}_M &= \frac{1}{8} \left| \begin{pmatrix} gW_3 + g'B & g(W_1 - iW_2) \\ g(W_1 + iW_2) & -gW_3 + g'B \end{pmatrix} \begin{pmatrix} 0 \\ v \end{pmatrix} \right|^2 \\ &= \frac{g^2 v^2}{8} \left[ W_1^2 + W_2^2 + \left( \frac{g'}{g} B - W_3 \right)^2 \right] \end{aligned} \quad (2.7)$$

Defining the Weinberg angle  $\tan(\theta_W) = g'/g$  and the following *physical* fields :

$$W^\pm = \frac{1}{\sqrt{2}} (W_1 \mp iW_2) \quad (2.8)$$

$$Z^0 = \cos \theta_W W_3 - \sin \theta_W B$$

$$A^0 = \sin \theta_W W_3 + \cos \theta_W B$$

we can write the piece of the Lagrangian associated to the vector boson masses as

$$\mathcal{L}_{MV} = \frac{1}{4}g^2 v^2 W^+ W^- + \frac{1}{8}(g^2 + g'^2)v^2 Z^0 Z^0. \quad (2.9)$$

and we have the following values of the masses for the vector bosons :

$$\begin{aligned} m_W^2 &= \frac{1}{4}v^2 g^2 \\ m_Z^2 &= \frac{1}{4}v^2(g^2 + g'^2) \\ m_A^2 &= 0 \end{aligned} \quad (2.10)$$

We thus see how the Higgs mechanism gives rise to the masses of the  $W^\pm$  and  $Z$  boson in the Standard Model; the mass of the photon is zero, as expected. The  $SU(2)_L \otimes U(1)_Y$  symmetry of the initially massless  $W_{1,2,3}$  and  $B$  fields is broken to the  $U(1)_{EM}$ . Of the four degrees of freedom in the complex Higgs doublet, three are “eaten” when we give mass to the  $W^\pm$  and  $Z_0$ , while the other degree of freedom is the Higgs particle, as found in 2012 by the ATLAS and CMS collaborations [51, 52].

## 181 Quantum Chromodynamics

Quantum chromodynamics (or the theory of the *strong* force) characterizes the behavior of *colored* particles, collectively known as *partons*. The partons of the Standard Model are the (fermionic) quarks, and the (bosonic) gluons. The strong force is governed by  $SU(3)_C$ , an unbroken symmetry in the Standard Model, which implies the gluon remains massless. Defining the covariant derivative for QCD as

$$D^\mu = \partial^\mu + ig_s G_a^\mu L_a, a = 1, \dots, 8 \quad (2.11)$$

where  $L_a$  are the generators of  $SU(3)_C$ , and  $g_s$  is the coupling constant of the strong force. The QCD Lagrangian then is given by

$$\mathcal{L}_{QCD} = i\bar{\psi}_f D_\mu \gamma^\mu \psi_f - \frac{1}{4}G_{a,\mu\nu} G_a^{\mu\nu} \quad (2.12)$$

where the summation over  $f$  is for quarks *families*, and  $G_a^{\mu\nu}$  is the gluon field strength tensor, given by

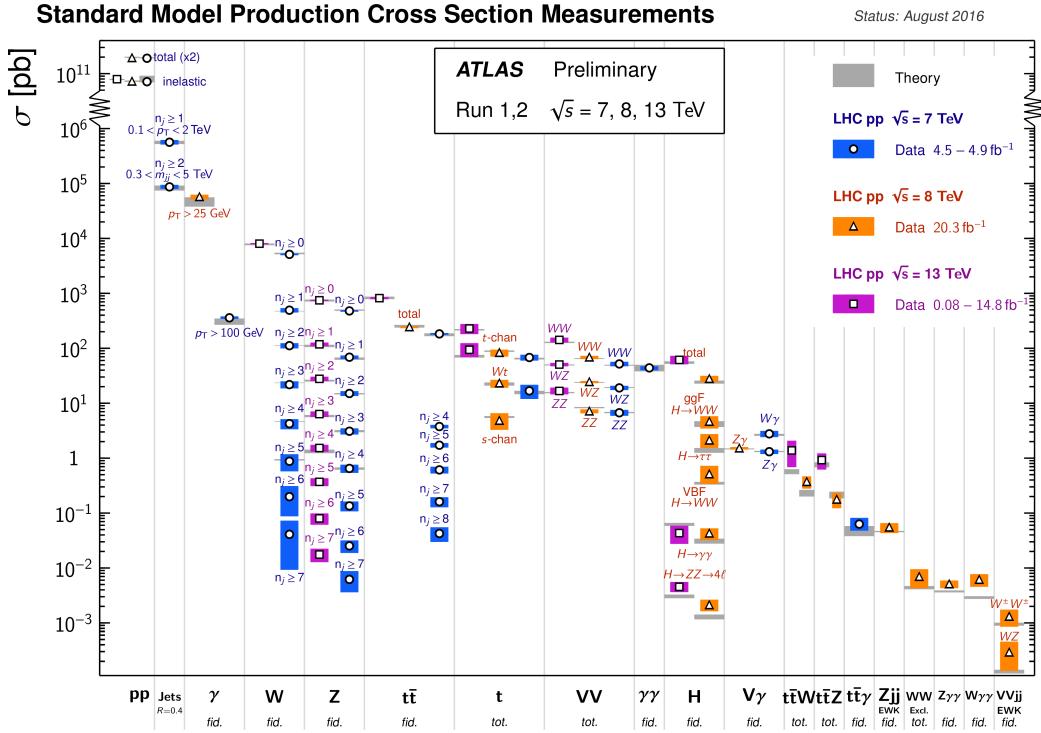
$$G_a^{\mu\nu} = \partial^\mu G_a^\nu - \partial^\nu G_a^\mu - g_s f^{abc} G_b^\mu G_c^\nu, a, b, c = 1, \dots, 8 \quad (2.13)$$

182 where  $f^{abc}$  are the structure constants of  $SU(3)_C$ , which are analogous to  $\epsilon_{abc}$  for  
 183  $SU(2)_L$ . The kinetic term for the quarks is contained in the standard  $\partial_\mu$  term, while  
 184 the field strength term contains the interactions between the quarks and gluons, as  
 185 well as the gluon self-interactions.

186 Written down in this simple form, the QCD Lagrangian does not seem much  
 187 different from the QED Lagrangian, with the proper adjustments for the different  
 188 group structures. The gluon is massless, like the photon, so one could naïvely expect  
 189 an infinite range force, and it pays to understand why this is not the case. The  
 190 reason for this fundamental difference is the gluon self-interactions arising in the  
 191 field strength tensor term of the Lagrangian. This leads to the phenomena of *color*  
 192 *confinement*, which describes how one only observes color-neutral particles alone in  
 193 nature. In contrast to the electromagnetic force, particles which interact via the  
 194 strong force experience a *greater* force as the distance between the particles increases.  
 195 At long distances, the potential is given by  $V(r) = -kr$ . At some point, it is more  
 196 energetically favorable to create additional partons out of the vacuum than continue  
 197 pulling apart the existing partons, and the colored particles undergo *fragmentation*.  
 198 This leads to *hadronization*. Bare quarks and gluons are actually observed as sprays  
 199 of hadrons (primarily kaons and pions); these sprays are known as *jets*, which are  
 200 what are observed by experiments.

201 It is important to recognize the importance of understanding these QCD inter-  
 202 actions in high-energy hadron colliders such as the LHC. Since protons are hadrons,  
 203 proton-proton collisions such as those produced by the LHC are primarily governed by  
 204 the processes of QCD. In particular, by far the most frequent process observed in LHC  
 205 experiments is dijet production from gluon-gluon interactions (see Fig.2.2). These

Figure 2.2: Cross-sections of various Standard Model processes



206 gluons that interact are part of the *sea* particles inside the proton; the simple  $p = uud$   
207 model does not apply. The main *valence*  $uud$  quarks are constantly interacting via  
208 gluons, which can themselves radiate gluons or split into quarks, and so on. A more  
209 useful understanding is given by the colloquially-known *bag* model [53, 54], where the  
210 proton is seen as a “bag” of (in principle) infinitely many partons, each with energy  
211  $E < \sqrt{s} = 6.5$  TeV. One then collides this (proton) bag with another, and views the  
212 products of this very complicated collision, where calculations include many loops in  
213 nonperturbative QCD calculations.

214 Fortunately, we are generally saved by the QCD factorization theorems [55]. This  
215 allows one to understand the hard (i.e. short distance or high energy)  $2 \rightarrow 2$  parton  
216 process using the tools of perturbative QCD, while making series of approximations  
217 known as a *parton shower* model to understand the additional corrections from  
218 nonperturbative QCD. We will discuss the reconstruction of jets by experiments in  
219 Ch.5.

220 **Fermions**

221 We will now look more closely at the fermions in the Standard Model [56].

222 As noted earlier in Sec.2.2, the fermions of the Standard Model can be first  
 223 distinguished between those that interact via the strong force (quarks) and those  
 224 which do not (leptons).

There are six leptons in the Standard Model, which can be placed into three  
*generations*.

$$\begin{pmatrix} e \\ \nu_e \end{pmatrix}, \begin{pmatrix} \mu \\ \nu_\mu \end{pmatrix}, \begin{pmatrix} \tau \\ \nu_\tau \end{pmatrix} \quad (2.14)$$

225 There is the electron ( $e$ ), muon ( $\mu$ ), and tau ( $\tau$ ), each of which has an associated  
 226 neutrino ( $\nu_e, \nu_\mu, \nu_\tau$ ). Each of the so-called charged (“electron-like”) leptons has  
 227 electromagnetic charge  $-1$ , while the neutrinos all have  $q_{EM} = 0$ .

228 Often in an experimental context, lepton is used to denote the stable electron  
 229 and metastable muon, due to their striking experimental signatures. Taus are often  
 230 treated separately, due to their much shorter lifetime of  $\tau_\tau \sim 10^{-13}s$ ; these decay  
 231 through hadrons or the other leptons, so often physics analyses at the LHC treat  
 232 them as jets or leptons, as will be done in this thesis.

233 As the neutrinos are electrically neutral, nearly massless, and only interact via the  
 234 weak force, it is quite difficult to observe them directly. Since LHC experiments rely  
 235 overwhelmingly on electromagnetic interactions to observe particles, the presence of  
 236 neutrinos is not observed directly. Neutrinos are instead observed by the conservation  
 237 of four-momentum in the plane transverse to the proton-proton collisions, known as  
 238 *missing transverse energy*.

There are six quarks in the Standard Model : up, down, charm, strange, top, and  
 bottom. Quarks are similar organized into three generations :

$$\begin{pmatrix} u \\ d \end{pmatrix}, \begin{pmatrix} c \\ s \end{pmatrix}, \begin{pmatrix} t \\ b \end{pmatrix} \quad (2.15)$$

239 where we speak of “up-like” quarks and “down-like” quarks.

240 Each up-like quark has charge  $q_{up} = 2/3$ , while the down-like quarks have  $q_{down} =$   
241  $-1/3$ . At the high energies of the LHC, one often makes the distinction between  
242 the light quarks ( $u, d, c, s$ ), the bottom quark, and top quark. In general, due to  
243 the hadronization process described above, the light quarks, with masses  $m_q < \sim$   
244  $1.5\text{GeV}$  are indistinguishable by LHC experiments. Their hadronic decay products  
245 generally have long lifetimes and they are reconstructed as jets.<sup>1</sup>. The bottom quark  
246 hadronizes primarily through the  $B$ -mesons, which generally travels a short distance  
247 before decaying to other hadrons. This allows one to distinguish decays via  $b$ -quarks  
248 from other jets; this procedure is known as *b-tagging* and will be discussed more in  
249 Ch.5. Due to its large mass, the top quark decays before it can hadronize; there  
250 are no bound states associated to the top quark. The top is of particular interest at  
251 the LHC; it has a striking signature through its most common decay mode  $t \rightarrow Wb$ .  
252 Decays via tops, especially  $t\bar{t}$  are frequently an important signal decay mode, or an  
253 important background process.

## 254 **Interactions in the Standard Model**

255 We briefly overview the entirety of the fundamental interactions of the Standard  
256 Model; these can also be found in 2.3.

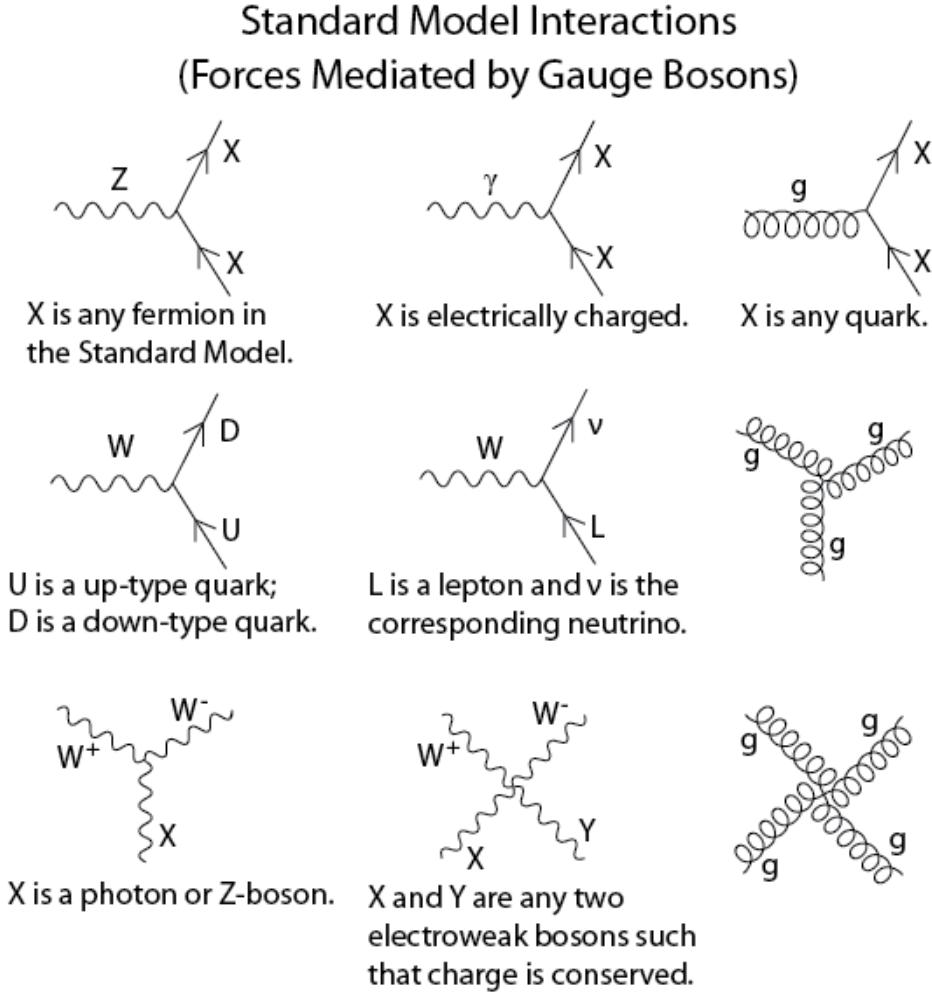
257 The electromagnetic force, mediated by the photon, interacts with via a three-  
258 point coupling all charged particles in the Standard Model. The photon thus interacts  
259 with all the quarks, the charged leptons, and the charged  $W^\pm$  bosons.

260 The weak force is mediated by three particles : the  $W^\pm$  and the  $Z^0$ . The  $Z^0$  can  
261 interacts with all fermions via a three-point coupling. A real  $Z_0$  can thus decay to  
262 a fermion-antifermion pair of all SM fermions except the top quark, due to its large

---

<sup>1</sup>In some contexts, charm quarks are also treated as a separate category, although it is quite difficult to distinguish charm quarks from the other light quarks.

Figure 2.3: The interactions of the Standard Model



mass. The  $W^\pm$  has two important three-point interactions with fermions. First, the  $W^\pm$  can interact with an up-like quark and a down-like quark; an important example in LHC experiments is  $t \rightarrow W b$ . The coupling constants for these interactions are encoded in the unitary matrix known as the Cabibbo–Kobayashi–Maskawa (CKM) matrix [57, 58], and are generally known as flavor-changing interactions. Secondly, the  $W^\pm$  interacts with a charged lepton and its corresponding neutrino. In this case, the unitary matrix that corresponds to CKM matrix for quarks is the identity matrix, which forbids (fundamental) vertices such as  $\mu \rightarrow We$ . For leptons, instead this is a two-step process :  $\mu \rightarrow \nu_m u W \rightarrow \nu_m u \bar{\nu}_e e$ . Finally, there are the self-interactions

272 of the weak gauge bosons. There is a three-point and four-point interaction; all  
273 combinations are allowed which conserve electric charge.

274 The strong force is mediated by the gluon, which as discussed above also carries  
275 the strong color charge. There is the fundamental three-point interaction, where a  
276 quark radiates a gluon. Additionally, there are the three-point and four-point gluon-  
277 only interactions.

## 278 2.3 Deficiencies of the Standard Model

279 At this point, it is quite easy to simply rest on our laurels. This relatively simple  
280 theory is capable of explaining a very wide range of phenomena, which ultimately  
281 break down only to combinations of nine diagrams shown in Fig.2.3. Unfortunately,  
282 there are some unexplained problems with the Standard Model. We cannot go  
283 through all of the potential issues in this thesis, but we will motivate the primary  
284 issues which naturally lead one to *supersymmetry*, as we will see in Ch.3.

The Standard Model has many free parameters; see Table 2.1 In general, we prefer models with less free parameters. A great example of this fact, and the primary experimental evidence for EWSB, is the relationship between the couplings of the weak force and the masses of the gauge bosons of the weak force :

$$\rho \equiv \frac{m_W^2}{m_Z^2 \cos^2 \theta_W} \stackrel{?}{=} 1 \quad (2.16)$$

285 where ? indicates that this is a testable prediction of the Standard Model (in  
286 particular, that the gauge bosons gain mass through EWSB). This relationship has  
287 been measured within experimental and theoretical predictions. We would like to  
288 produce additional such relationships, which would exist if the Standard Model is a  
289 low-energy approximation of some other theory.

290 An additional issue is the lack of *gauge coupling unification*. The couplings of  
291 any quantum field theory “run” as a function of the distance scales (or inversely,

Table 2.1: Parameters of the Standard Model. For values dependent on the renormalization scheme, we use a combination of the on-shell normalization scheme [59–62] and modified minimal subtraction scheme with  $m_{\bar{MS}}$  as indicated in the table[63]

$m_e$	Electron mass	511 keV
$m_\mu$	Muon mass	105.7 MeV
$m_\tau$	Tau mass	1.78 GeV
$m_u$	Up quark mass	1.9 MeV ( $m_{\bar{MS}} = 2\text{GeV}$ )
$m_d$	Down quark mass	4.4 MeV ( $m_{\bar{MS}} = 2\text{GeV}$ )
$m_s$	Strange quark mass	87 MeV ( $m_{\bar{MS}} = 2\text{GeV}$ )
$m_c$	Charm quark mass	1.32 GeV ( $m_{\bar{MS}} = m_c$ )
$m_b$	Bottom quark mass	4.24 GeV ( $m_{\bar{MS}} = m_b$ )
$m_t$	Top quark mass	172.7 GeV (on-shell renormalization)
$\theta_{12}$ CKM	12-mixing angle	13.1°
$\theta_{23}$ CKM	23-mixing angle	2.4°
$\theta_{13}$ CKM	13-mixing angle	0.2°
$\delta$ CKM	CP-violating Phase	0.995
$g'$	U(1) gauge coupling	0.357 ( $m_{\bar{MS}} = m_Z$ )
$g$	SU(2) gauge coupling	0.652 ( $m_{\bar{MS}} = m_Z$ )
$g_s$	SU(3) gauge coupling	1.221 ( $m_{\bar{MS}} = m_Z$ )
$\theta_{QCD}$	QCD vacuum angle	~0
VEV	Higgs vacuum expectation value	246 GeV
$m_H$	Higgs mass	125 GeV

292 energy scales) of the theory. The idea is closely related to the unification of the  
 293 electromagnetic and weak forces at the so-called *electroweak scale* of  $O(100 \text{ GeV})$ .

294 One would hope this behavior was repeated between the electroweak forces and the  
 295 strong force at some suitable energy scale. The Standard Model does automatically  
 296 not exhibit this behavior, as we can see in Fig.2.4.

The most significant problem with the Standard Model is the *hierarchy problem*. In its most straightforward incarnation, the Higgs scalar field is subject to quantum corrections through loop diagrams, as shown in Fig.2.5. For demonstration, we use the contributions from the top quark, since the top quark has the largest Higgs Yukawa coupling due to its large mass. In general, we should expect these corrections to quadratically depend on the scale of the ultraviolet physics,  $\Lambda$ . Briefly assume there is no new physics before the Planck scale of gravity,  $\Lambda_{\text{Planck}} = 10^{19} \text{ GeV}$ . In this

Figure 2.4: The running of Standard Model gauge couplings. The Standard Model couplings do not unify at high energies, which indicates it cannot completely describe nature through the Planck scale.

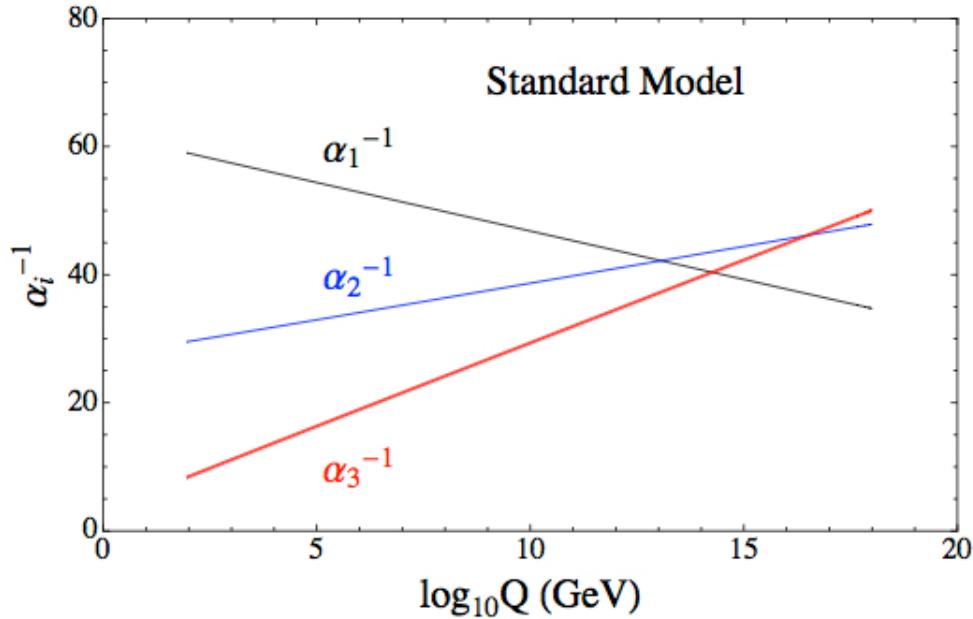
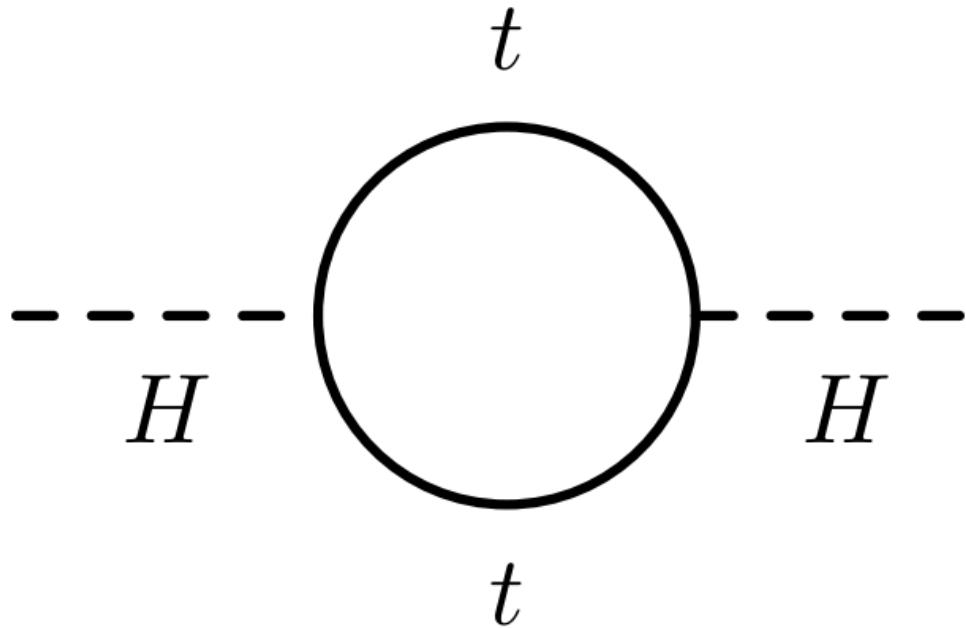


Figure 2.5: The dominant quantum loop correction to the Higgs mass in the Standard Model.



case, we expect the corrections to the Higgs mass like

$$\delta m_H^2 \approx \left( \frac{m_t}{8\pi^2 \langle \phi \rangle_{VEV}} \right)^2 \Lambda_{Planck}^2. \quad (2.17)$$

297 To achieve the miraculous cancellation required to get the observed Higgs mass of  
298 125 GeV, one needs to then set the bare Higgs mass  $m_0$ , our input to the Standard  
299 Model Lagrangian, itself to a *precise* value  $\sim 10^{19}$  GeV. This extraordinary level of  
300 parameter finetuning is quite undesirable, and within the framework of the Standard  
301 Model, there is little that can be done to alleviate this issue.

302 An additional concern, of a different nature, is the lack of a *dark matter* candidate  
303 in the Standard Model. Dark matter was discovered by observing galactic rotation  
304 curves, which showed that much of the matter that interacted gravitationally was  
305 invisible to our (electromagnetic) telescopes [16–22]. The postulation of the existence  
306 of dark matter, which interacts at least through gravity, allows one to understand  
307 these galactic rotation curves. Unfortunately, no particle in the Standard Model could  
308 possibly be the dark matter particle. The only candidate truly worth another look is  
309 the neutrino, but it has been shown that the neutrino content of the universe is simply  
310 too small to explain the galactic rotation curves [22, 64]. The experimental evidence  
311 from the galactic rotations curves thus show there *must* be additional physics beyond  
312 the Standard Model, which is yet to be understood.

313 In the next chapter, we will see how these problems can be alleviated by the theory  
314 of supersymmetry.

Figure 2.6: Particles of the Standard Model

mass → $\approx 2.3 \text{ MeV}/c^2$ charge → $2/3$ spin → $1/2$	mass → $\approx 1.275 \text{ GeV}/c^2$ charge → $2/3$ spin → $1/2$	mass → $\approx 173.07 \text{ GeV}/c^2$ charge → $2/3$ spin → $1/2$	mass → $0$ charge → $0$ spin → $1$	mass → $\approx 126 \text{ GeV}/c^2$ charge → $0$ spin → $0$
<b>QUARKS</b> <b>u</b> up	<b>c</b> charm	<b>t</b> top	<b>g</b> gluon	<b>H</b> Higgs boson
$\approx 4.8 \text{ MeV}/c^2$ $-1/3$ $1/2$ down	$\approx 95 \text{ MeV}/c^2$ $-1/3$ $1/2$ strange	$\approx 4.18 \text{ GeV}/c^2$ $-1/3$ $1/2$ bottom	$\gamma$	photon
<b>LEPTONS</b> <b>e</b> electron	<b><math>\mu</math></b> muon	<b><math>\tau</math></b> tau	<b>Z</b> Z boson	<b>Gauge Bosons</b>
$<2.2 \text{ eV}/c^2$ $0$ $1/2$ electron neutrino	$<0.17 \text{ MeV}/c^2$ $0$ $1/2$ muon neutrino	$<15.5 \text{ MeV}/c^2$ $0$ $1/2$ tau neutrino	$80.4 \text{ GeV}/c^2$ $\pm 1$ $1$ W boson	



*Supersymmetry*

317 This chapter will introduce supersymmetry (SUSY) [15, 23–35]. We will begin by  
 318 introducing the concept of a *superspace*, and discuss some general ingredients of  
 319 supersymmetric theories. This will include a discussion of how the problems with the  
 320 Standard Model described in Ch.2 are naturally fixed by these theories.

321 The next step is to discuss the particle content of the *Minimally Supersymmetric*  
 322 *Standard Model* (MSSM). As its name implies, this theory contains the minimal  
 323 additional particle content to make Standard Model supersymmetric. We then discuss  
 324 the important phenomenological consequences of this theory, especially as it would  
 325 be observed in experiments at the LHC.

326 **3.1 Supersymmetric theories : from space to  
 327 superspace**

328 **Coleman-Mandula “no-go” theorem**

329 We begin the theoretical motivation for supersymmetry by citing the “no-go” theorem  
 330 of Coleman and Mandula [65]. This theorem forbids *spin-charge unification*; it  
 331 states that all quantum field theories which contain nontrivial interactions must be  
 332 a direct product of the Poincaré group of Lorentz symmetries, the internal product  
 333 from of gauge symmetries, and the discrete symmetries of parity, charge conjugation,  
 334 and time reversal. The assumptions which go into building the Coleman-Mandula

theorem are quite restrictive, but there is one unique way out, which has become known as *supersymmetry* [26, 66]. In particular, we must introduce a *spinorial* group generator  $Q$ . Alternatively, and equivalently, this can be viewed as the addition of anti-commuting coordinates; space plus these new anti-commuting coordinates is then called *superspace* [67]. We will not investiage this view in detail, but it is also a quite intuitive and beautiful way to construct supersymmetry[15].

## 341 Supersymmetry transformations

342 A *supersymmetric* transformation  $Q$  transforms a bosonic state into a fermionic state,  
 343 and vice versa :

$$Q |\text{Fermion}\rangle = |\text{Boson}\rangle \quad (3.1)$$

$$Q |\text{Boson}\rangle = |\text{Fermion}\rangle \quad (3.2)$$

To ensure this relation holds,  $Q$  must be an anticommuting spinor. Additionally, since spinors are inherently complex,  $Q^\dagger$  must also be a generator of the supersymmetry transformation. Since  $Q$  and  $Q^\dagger$  are spinor objects (with  $s = 1/2$ ), we can see that supersymmetry must be a spacetime symmetry. The Haag-Lopuszanski-Sohnius extension [66] of the Coleman-Mandula theorem [65] is quite restrictive about the forms of such a symmetry. Here, we simply write the (anti-) commutation relations [15] :

$$Q_\alpha, Q_{\dot{\alpha}}^\dagger = -2\sigma_{\alpha\dot{\alpha}\mu} P_\mu \quad (3.3)$$

$$Q_\alpha, Q_{\dot{\beta}} = Q_{\dot{\alpha}}^\dagger, Q_{\dot{\beta}}^\dagger = 0 \quad (3.4)$$

$$[P^\mu, Q_\alpha] = [P^\mu, Q_{\dot{\alpha}}^\dagger] = 0 \quad (3.5)$$

<sup>344</sup> **Supermultiplets**

<sup>345</sup> In a supersymmetric theory, we organize single-particle states into irreducible  
<sup>346</sup> representations of the supersymmetric algebra which are known as *supermultiplets*.  
<sup>347</sup> Each supermultiplet contains a fermion state  $|F\rangle$  and a boson state  $|B\rangle$ ; these two  
<sup>348</sup> states are the known as *superpartners*. These are related by some combination of  
<sup>349</sup>  $Q$  and  $Q^\dagger$ , up to a spacetime transformation.  $Q$  and  $Q^\dagger$  commute with the mass-  
<sup>350</sup> squared operator  $-P^2$  and the operators corresponding to the gauge transformations  
<sup>351</sup> [15]; in particular, the gauge interactions of the Standard Model. In an unbroken  
<sup>352</sup> supersymmetric theory, this means the states  $|F\rangle$  and  $|B\rangle$  have exactly the same mass,  
<sup>353</sup> electromagnetic charge, electroweak isospin, and color charges. One can also prove  
<sup>354</sup> [15] that each supermultiplet contains the exact same number of bosonic ( $n_B$ ) and  
<sup>355</sup> fermion ( $n_F$ ) degrees of freedom. We now explore the possible types of supermultiples  
<sup>356</sup> one can find in a renormalizable supersymmetric theory.

<sup>357</sup> Since each supermultiplet must contain a fermion state, the simplest type of  
<sup>358</sup> supermultiplet contains a single Weyl fermion state ( $n_F = 2$ ) which is paired with  
<sup>359</sup>  $n_B = 2$  scalar bosonic degrees of freedom. This is most conveniently constructed as  
<sup>360</sup> single complex scalar field. We call this construction a *scalar supermultiplet* or *chiral*  
<sup>361</sup> *supermultiplet*. The second name is indicative; only chiral supermultiplets can contain  
<sup>362</sup> fermions whose right-handed and left-handed components transform differently under  
<sup>363</sup> the gauge interactions (as of course happens in the Standard Model).

<sup>364</sup> The second type of supermultiplet we construct is known as a *gauge* supermul-  
<sup>365</sup> tiplet. We take a spin-1 gauge boson (which must be massless due to the gauge  
<sup>366</sup> symmetry, so  $n_B = 2$ ) and pair this with a single massless Weyl spinor<sup>1</sup>. The gauge  
<sup>367</sup> bosons transform as the adjoint representation of the their respective gauge groups;  
<sup>368</sup> their fermionic partners, which are known as gauginos, must also. In particular,  
<sup>369</sup> the left-handed and right-handed components of the gaugino fermions have the same

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<sup>1</sup>Choosing an  $s = 3/2$  massless fermion leads to nonrenormalizable interactions.

370 gauge transformation properties.

371 Excluding gravity, this is the entire list of supermultiplets which can participate  
372 in renormalizable interactions in what is known as  $N = 1$  supersymmetry. This  
373 means there is only one copy of the supersymmetry generators  $Q$  and  $Q^\dagger$ . This is  
374 essentially the only “easy” phenomenological choice, since it is the only choice in four  
375 dimensions which allows for the chiral fermions and parity violations built into the  
376 Standard Model, and we will not look further into  $N > 1$  supersymmetry in this thesis.

377 The primary goal, after understanding the possible structures of the multiplets  
378 above, is to fit the Standard Model particles into a multiplet, and therefore make  
379 predictions about their supersymmetric partners. We explore this in the next section.

## 380 3.2 Minimally Supersymmetric Standard Model

381 To construct what is known as the MSSM [susyPrimer , 68–71], we need a few  
382 ingredients and assumptions. First, we match the Standard Model particles with  
383 their corresponding superpartners of the MSSM. We will also introduce the naming  
384 of the superpartners (also known as *sparticles*). We discuss a very common additional  
385 restraint imposed on the MSSM, known as  $R$ –parity. We also discuss the concept of  
386 soft supersymmetry breaking and how it manifests itself in the MSSM.

### 387 Chiral supermultiplets

388 The first thing we deduce is directly from Sec.?? . The bosonic superpartners  
389 associated to the quarks and leptons *must* be spin 0, since the quarks and leptons must  
390 be arranged in a chiral supermultiplet. This is essentially the note above, since the  
391 chiral supermultiplet is the only one which can distinguish between the left-handed  
392 and right-handed components of the Standard Model particles. The superpartners of  
393 the quarks and leptons are known as *squarks* and *sleptons*, or *sfermions* in aggregate.

394 (for ‘‘scalar quarks’’, ‘‘scalar leptons’’, and ‘‘scalar fermion’’<sup>2</sup>). The ‘‘s-’’ prefix  
 395 can also be added to the individual quarks i.e. *selectron*, *sneutrino*, and *stop*. The  
 396 notation is to add a  $\sim$  over the corresponding Standard Model particle i.e.  $\tilde{e}$ , the  
 397 selectron is the superpartner of the electron. The two-component Weyl spinors of the  
 398 Standard Model must each have their own (complex scalar) partner i.e.  $e_L, e_R$  have  
 399 two distinct partners :  $\tilde{e}_L, \tilde{e}_R$ . As noted above, the gauge interactions of any of the  
 400 sfermions are identical to those of their Standard Model partners.

Due to the scalar nature of the Higgs, it must obviously lie in a chiral supermultiplet. To avoid gauge anomalies and ensure the correct Yukawa couplings to the quarks and leptons[15], we must add additional Higgs bosons to any supersymmetric theory. In the MSSM, we have two chiral supermultiplets. The SM (SUSY) parts of the multiplets are denoted  $H_u(\tilde{H}_u)$  and  $H_d(\tilde{H}_d)$ . Writing out  $H_u$  and  $H_d$  explicitly:

$$H_u = \begin{pmatrix} H_u^+ \\ H_u^0 \end{pmatrix} \quad (3.6)$$

$$H_d = \begin{pmatrix} H_d^0 \\ H_d^- \end{pmatrix} \quad (3.7)$$

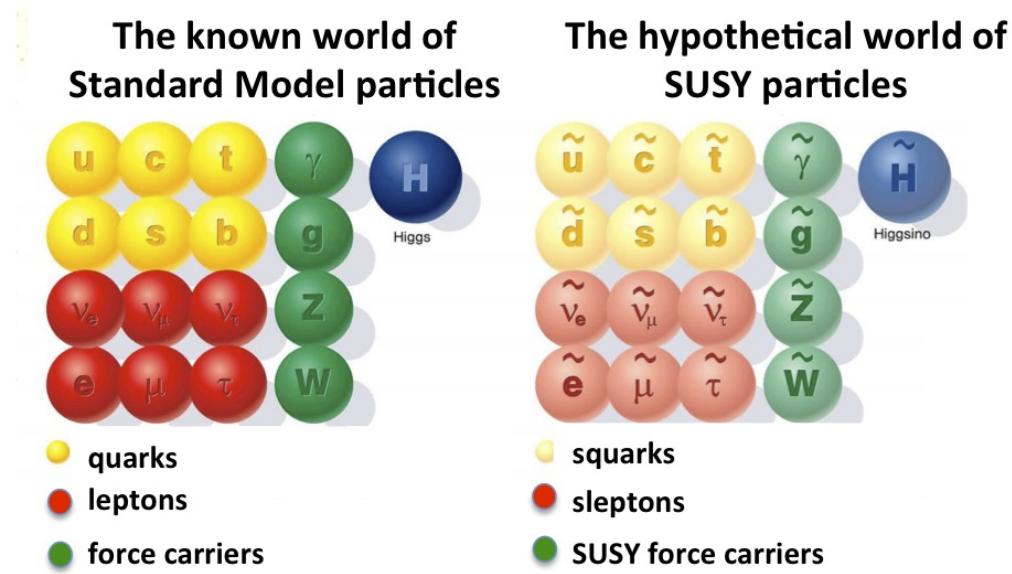
(3.8)

401 we see that  $H_u$  looks very similar to the SM Higgs with  $Y = 1$ , and  $H_d$  is symmetric  
 402 to this with  $+ \rightarrow -$ , with  $Y = -1$ . The SM Higgs boson,  $h_0$ , is a linear superposition  
 403 of the neutral components of these two doublets. The SUSY parts of the Higgs  
 404 multiplets,  $\tilde{H}_u$  and  $\tilde{H}_d$ , are each left-handed Weyl spinors. For generic spin-1/2  
 405 sparticles, we add the ‘‘-ino’’ suffix. We then call the partners of the two Higgs  
 406 collectively the *Higgsinos*.

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<sup>2</sup>The last one should probably have bigger scare quotes.

Figure 3.1: Particles of the MSSM



## 407 Gauge supermultiplets

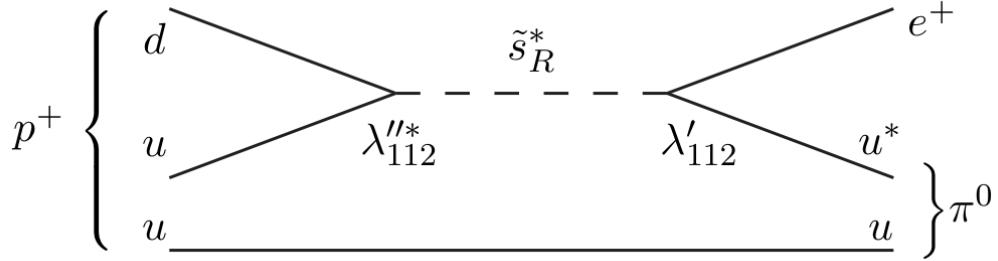
408 The superpartners of the gauge bosons must all be in gauge supermultiplets since  
 409 they contain a spin-1 particle. Collectively, we refer to the superpartners of the  
 410 gauge bosons as the gauginos.

411 The first gauge supermultiplet contains the gluon, and its superpartner, which is  
 412 known as the *gluino*, denoted  $\tilde{g}$ . The gluon is of course the SM mediator of  $SU(3)_C$ ;  
 413 the gluino is also a colored particle, subject to  $SU(3)_C$ . From the SM before EWSB,  
 414 we have the four gauge bosons of the electroweak symmetry group  $SU(2)_L \otimes U(1)_Y$  :  
 415  $W^{1,2,3}$  and  $B^0$ . The superpartners of these particles are thus the *winos*  $W^{\tilde{1},\tilde{2},\tilde{3}}$  and  
 416 *bino*  $\tilde{B}^0$ , where each is placed in another gauge supermultiplet with its corresponding  
 417 SM particle. After EWSB, without breaking supersymmetry, we would also have the  
 418 zino  $\tilde{Z}^0$  and photino  $\tilde{\gamma}$ .

419 The entire particle content of the MSSM can be seen in Fig.3.1.

420 At this point, it's important to take a step back. Where are these particles?  
 421 As stated above, supersymmetric theories require that the masses and all quantum

Figure 3.2: This Feynmann diagram shows how proton decay is induced in the MSSM, if one does not impose  $R$ -parity.



422 numbers of the SM particle and its corresponding sparticle are the same. Of course,  
 423 we have not observed a selectron, squark, or wino. The answer, as it often is, is that  
 424 supersymmetry is *broken* by the vacuum state of nature [15].

## 425 **$R$ -parity**

This section is a quick aside to the general story.  $R$  – parity refers to an additional discrete symmetry which is often imposed on supersymmetric models. For a given particle state, we define

$$R = (-1)^{3(B-L)+2s} \quad (3.9)$$

426 where  $B, L$  is the baryon (lepton) number and  $s$  is the spin. The imposition of  
 427 this symmetry forbids certain terms from the MSSM Lagrangian that would violate  
 428 baryon and/or lepton number. This is required in order to prevent proton decay, as  
 429 shown in Fig.3.2<sup>3</sup>. .

430 In supersymmetric models, this is a  $\mathbb{Z}_2$  symmetry, where SM particles have  $R = 1$   
 431 and sparticles have  $R = -1$ . We will take  $R$  – parity as part of the definition of  
 432 the MSSM. We will discuss later the *drastic* consequences of this symmetry on SUSY  
 433 phenomenology

---

<sup>3</sup>Proton decay can actually be prevented by allowing only one of the four potential R-parity violating terms to survive.

434 **Soft supersymmetry breaking**

The fundamental idea of *soft* supersymmetry breaking[15, 34, 35, 72, 73] is that we would like to break supersymmetry without reintroducing the quadratic divergences we discussed at the end of Chapter 2. We write the Lagrangian in a form :

$$\mathcal{L}_{\text{MSSM}} = \mathcal{L}_{\text{SUSY}} + \mathcal{L}_{\text{soft}} \quad (3.10)$$

435 In this sense, the symmetry breaking is “soft”, since we have separated out the  
 436 completely symmetric terms from those soft terms which will not allow the quadratic  
 437 divergences to the Higgs mass.

438 The explicitly allowed terms in the soft-breaking Lagrangian are [35].

- 439 • Mass terms for the scalar components of the chiral supermultipletss  
 440 • Mass terms for the Weyl spinor components of the gauge supermultipletss  
 441 • Trilinear couplings of scalar components of chiral supermultiplets

In particular, using the field content described above for the MSSM, the softly-broken portion of the MSSM Lagrangian can be written

$$\mathcal{L}_{\text{soft}} = -\frac{1}{2} \left( M_3 \tilde{g} \tilde{g} + M_2 \tilde{W} \tilde{W} + M_1 \tilde{B} \tilde{B} + c.c. \right) \quad (3.11)$$

$$- \left( \tilde{u} a_u \tilde{Q} H_u - \tilde{d} a_d \tilde{Q} H_d - \tilde{e} a_e \tilde{L} H_d + c.c. \right) \quad (3.12)$$

$$- \tilde{Q}^\dagger m_Q^2 \tilde{Q} - \tilde{L}^\dagger m_L^2 \tilde{L} - \tilde{u} m_u^2 \tilde{u}^\dagger - \tilde{d} m_d^2 \tilde{d}^\dagger - \tilde{e} m_e^2 \tilde{e}^\dagger \quad (3.13)$$

$$- m_{H_u}^2 H_u^* H_u - m_{H_d}^2 H_d^* H_d - (b H_u H_d + cc). \quad (3.14)$$

442 where we have introduced the following notations :

443 1.  $M_3, M_2, M_1$  are the gluino, wino, and bino masses.

444 2.  $a_u, a_d, a_e$  are complex  $3 \times 3$  matrices in family space.

445 3.  $m_Q^2, m_u^2, m_d^2, m_L^2, m_e^2$  are hermitian  $3 \times 3$  matrices in family space.

446 4.  $m_{H_u}^2, m_{H_d}^2, b$  are the SUSY-breaking contributions to the Higgs potential.

447 We have written matrix terms without any sort of additional notational decoration  
 448 to indicate their matrix nature, and we now show why. The first term 1 are  
 449 straightforward; these are just the straightforward mass terms for these fields. There  
 450 are strong constraints on the off-diagonal terms for the matrices of 2 [74, 75]; for  
 451 simplicity, we will assume that each  $a_i, i = u, d, e$  is proportional to the Yukawa  
 452 coupling matrix :  $a_i = A_{i0}y_i$ . The matrices in ?? can be similarly constrained by  
 453 experiments [68, 75–82] Finally, we assume that the elements 4 contributing to the  
 454 Higgs potential as well as all of the 1 terms must be real, which limits the possible  
 455 CP-violating interactions to those of the Standard Model. We thus only consider  
 456 flavor-blind, CP-conserving interactions within the MSSM.

The important mixing for mass and gauge interaction eigenstates in the MSSM occurs within electroweak sector, in a process akin to EWSB in the Standard Model. The neutral portions of the Higgsinos doublets and the neutral gauginos ( $\tilde{H}_u^0, \tilde{H}_d^0, \tilde{B}^0, \tilde{W}^0$ ) of the gauge interaction basis mix to form what are known as the *neutralinos* of mass basis :

$$M_{\tilde{\chi}} = \begin{pmatrix} M_1 & 0 & -c_\beta s_W m_Z & s_\beta s_W m_Z \\ 0 & M_2 & c_\beta c_W m_Z & -s_\beta c_W m_Z \\ -c_\beta s_W m_Z & c_\beta c_W m_Z & 0 & -\mu \\ s_\beta s_W m_Z & -s_\beta c_W m_Z & -\mu & 0 \end{pmatrix} \quad (3.15)$$

457 where  $s(c)$  are the sine and cosine of angles related to EWSB, which introduced  
 458 masses to the gauginos and higgsinos. Diagonalization of this matrix gives the four  
 459 neutralino mass states, listed without loss of generality in order of increasing mass :  
 460  $\tilde{\chi}_{1,2,3,4}^0$ .

461 The neutralinos, especially the lightest neutralino  $\tilde{\chi}_1^0$ , are important ingredients  
 462 in SUSY phenomenology.

463     The same process can be done for the electrically charged gauginos with  
464     the charged portions of the Higgsino doublets along with the charged winos  
465      $(\tilde{H}_u^+, \tilde{H}_d^+, \tilde{W}^+, \tilde{W}^-)$ . This leads to the *charginos*, again in order of increasing mass  
466     :  $\tilde{\chi}_{1,2}^\pm$ .

467     

### 3.3 Phenomenology

468     We are finally at the point where we can discuss the phenomenology of the MSSM,  
469     in particular as it manifests itself at the energy scales of the LHC.

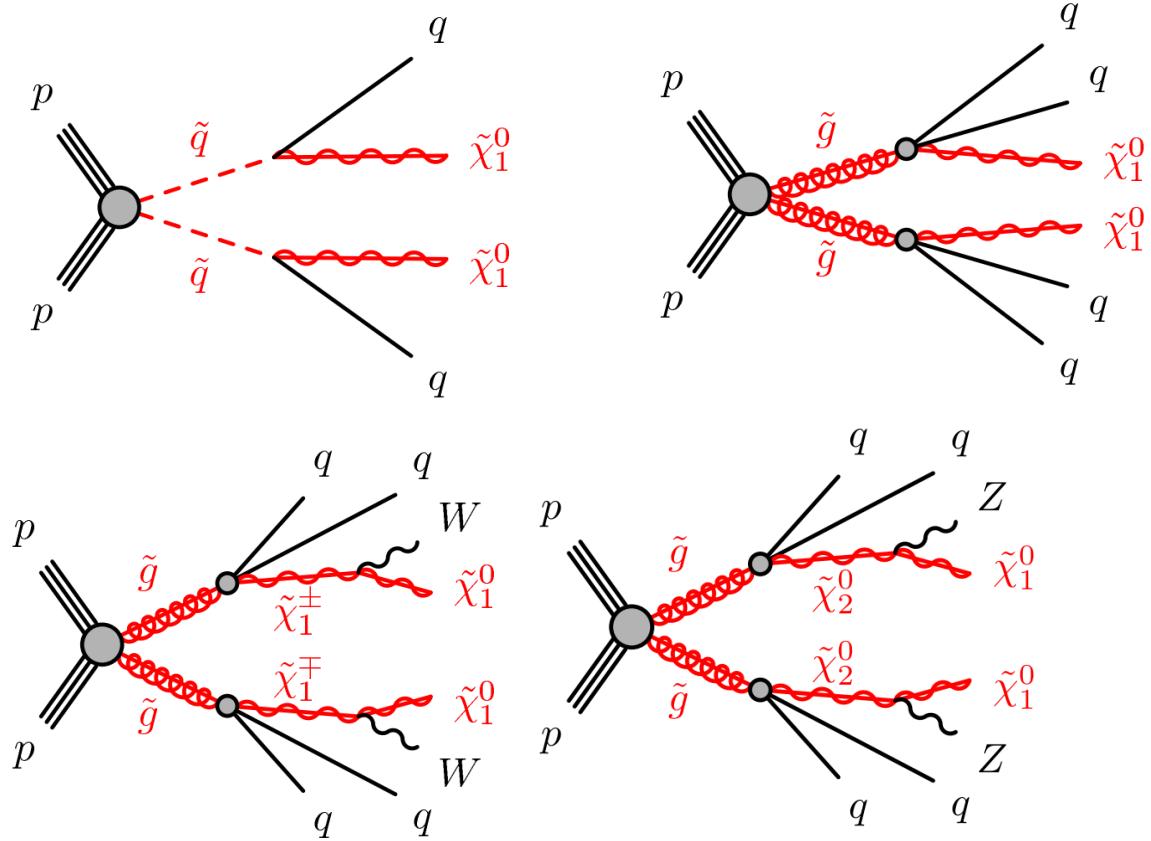
470     As noted above in Sec.3.2, the assumption of *R*–parity has important conse-  
471     quences for MSSM phenomenology. The SM particles have  $R = 1$ , while the sparticles  
472     all have  $R = -1$ . Simply, this is the “charge” of supersymmetry. Since the particles of  
473     LHC collisions ( $pp$ ) have total incoming  $R = 1$ , we must expect that all sparticles will  
474     be produced in *pairs*. An additional consequence of this symmetry is the fact that the  
475     lightest supersymmetric particle (LSP) is *stable*. Off each branch of the Feynmann  
476     diagram shown in Fig., we have  $R = -1$ , and this can only decay to another sparticle  
477     and a SM particle. Once we reach the lightest sparticle in the decay, it is absolutely  
478     stable. This leads to the common signature  $E_T^{\text{miss}}$  for a generic SUSY signal.

479     For this thesis, we will be presenting an inclusive search for squarks and gluinos  
480     with zero leptons in the final state. This is a very interesting decay channel<sup>4</sup>, due  
481     to the high cross-sections of  $\tilde{g}\tilde{g}$  and  $\tilde{q}\tilde{q}$  decays, as can be seen in Fig.?? [83]. This  
482     is a direct consequence of the fact that these are the colored particles of the MSSM.  
483     Since the sparticles interact with the gauge groups of the SM in the same way as their  
484     SM partners, the colored sparticles, the squarks and gluinos, are produced and decay  
485     as governed by the color group  $SU(3)_C$  with the strong coupling  $g_S$ . The digluino  
486     production is particularly copious, due to color factor corresponding to the color octet

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<sup>4</sup>Prior to Run1, probably the most *most* interesting SUSY decay channel.

Figure 3.3: SUSY signals considered in this thesis



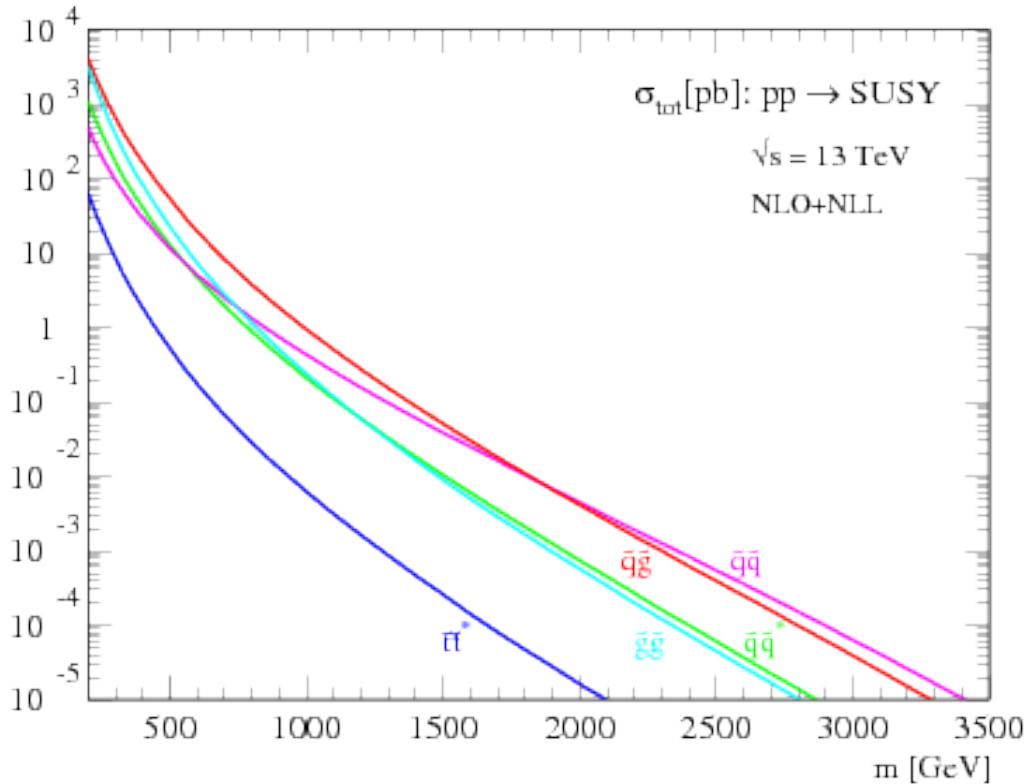
487 of  $SU(3)C$ .

488 In the case of disquark production, the most common decay mode of the squark in  
 489 the MSSM is a decay directly to the LSP plus a single SM quark [15]. This means the  
 490 basic search strategy of disquark production is two jets from the final state quarks,  
 491 plus missing transverse energy for the LSPs. There are also cascade decays, the most  
 492 common of which, and the only one considered in this thesis, is  $\tilde{q} \rightarrow q\tilde{\chi}_1^\pm \rightarrow qW^\pm\chi_1^0$ .

493 For digluino production, the most common decay is  $\tilde{g} \rightarrow g\tilde{q}$ , due to the large  
 494  $g_S$  coupling. The squark then decays as listed above. In this case, we generically  
 495 search for four jets and missing transverse energy from the LSPs. We can also have  
 496 the squark decay in association with a  $W^\pm$  or  $Z^0$ ; in this thesis, we are interested in  
 497 those cases where this vector boson goes hadronically.

498 In the context of experimental searches for SUSY, we often consider *simplified*

Figure 3.4: SUSY production cross-sections as a function of sparticle mass at  $\sqrt{s} = 13$  TeV.

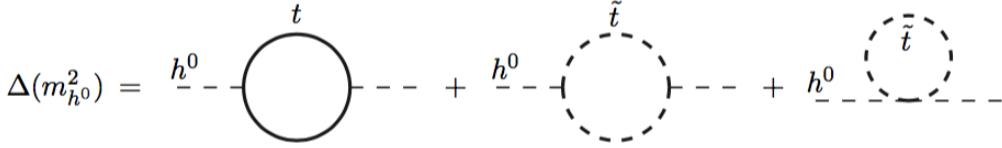


499 *models*. These models make certain assumptions which allow easy comparisons of  
 500 results by theorists and rival experimentalists. In the context of this thesis, the  
 501 simplified models will make assumptions about the branching ratios described in the  
 502 preceding paragraphs. In particular, we will often choose a model where the decay of  
 503 interest occurs with 100% branching ratio. This is entirely for ease of interpretation  
 504 by other physicists<sup>5</sup>, but it is important to recognize that these are more a useful  
 505 comparison tool, especially with limits, than a strict statement about the potential  
 506 masses of sought-after beyond the Standard Model particle.

---

<sup>5</sup>In the author's opinion, this often leads to more confusion than comprehension. We will revisit the shortcomings of simplified models in the Conclusion to this thesis.

Figure 3.5: Loop diagrams correct the Higgs mass in the MSSM



## 507 3.4 How SUSY solves the problems with the SM

508 We now return to the issues with the Standard Model as described in Ch.2 to see  
 509 how these issues are solved by supersymmetry.

### 510 Quadratic divergences to the Higgs mass

The quadratic divergences induced by the loop corrections to the Higgs mass, for example from the top Yukawa coupling, goes as

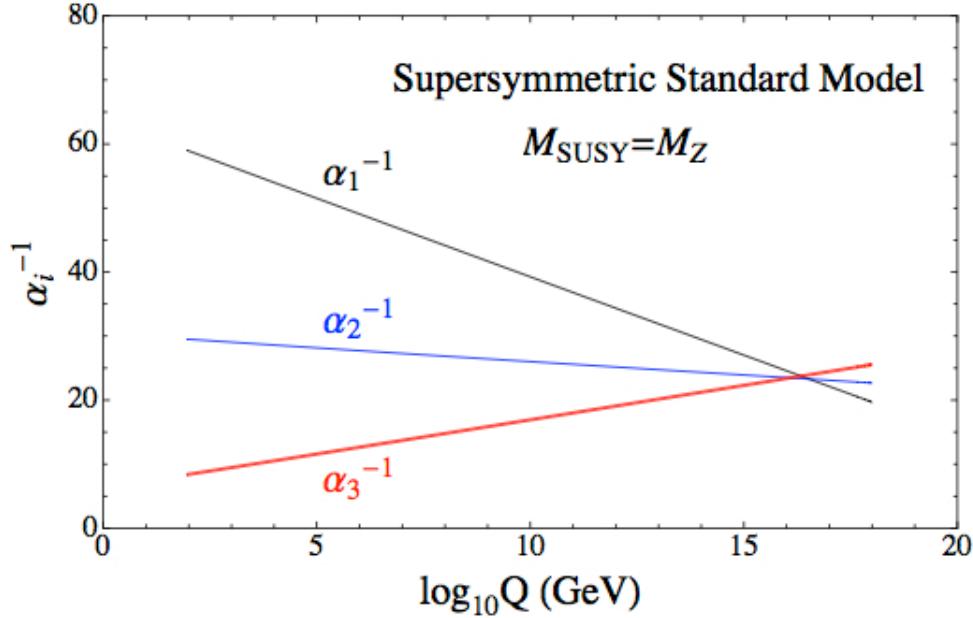
$$\delta m_H^2 \approx \left( \frac{m_t}{8\pi^2 \langle \phi \rangle_{VEV}} \right)^2 \Lambda_{Planck}^2. \quad (3.16)$$

511 The miraculous thing about SUSY is each of these terms *automatically* comes  
 512 with a term which exactly cancels this contribution[15]. The fermions and bosons  
 513 have opposite signs in this loop diagram to all orders in perturbation theory, which  
 514 completely solves the hierarchy problem. This is the most well-motivated reason for  
 515 supersymmetry.

### 516 Gauge coupling unification

517 An additional motivation for supersymmetry is seen by the gauge coupling unification  
 518 high scales. In the Standard Model, as we saw the gauge couplings fail to unify at  
 519 high energies. In the MSSM and many other forms of supersymmetry, the gauge  
 520 couplings unify at high energy, as can be seen in Fig.???. This provides additional  
 521 aesthetic motivation for supersymmetric theories.

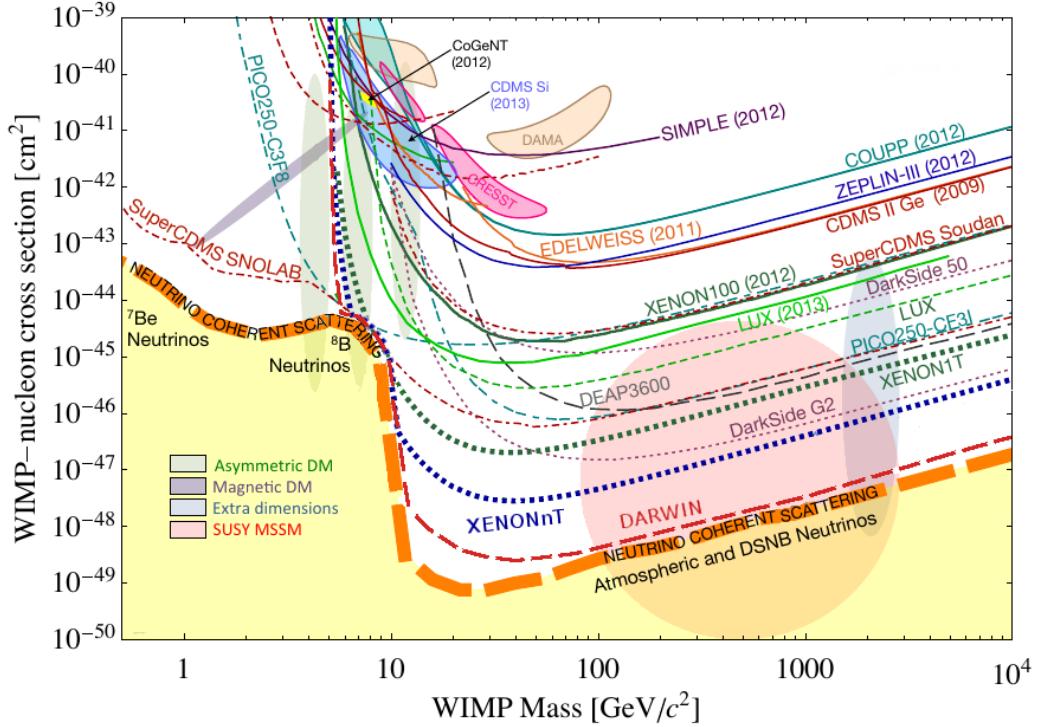
Figure 3.6: The running of Standard Model gauge couplings; compare to Fig.2.4. The MSSM gauge couplings nearly intersect at high energies.



## 522 Dark matter

523 As we discussed previously, the lack of any dark matter candidate in the Standard  
 524 Model naturally leads to beyond the Standard Model theories. In the Standard Model,  
 525 there is a natural dark matter candidate in the lightest supersymmetric particle[15]  
 526 The LSP would in dark matter experiments be called a *weakly-interacting massive*  
 527 *particle* (WIMP), which is a type of cold dark matter [22, 84]. These WIMPS would  
 528 only interact through the weak force and gravity, which is exactly as a model like the  
 529 MSSM predicts for the neutralino. In Fig.3.7, we can see the current WIMP exclusions  
 530 for a given mass. The range of allowed masses which have not been excluded for LSPs  
 531 and WIMPs have significant overlap. This provides additional motivation outside of  
 532 the context of theoretical details.

Figure 3.7: WIMP exclusions from direct dark matter detection experiments.



## 533 3.5 Conclusions

534 Supersymmetry is the most well-motivated theory for physics beyond the Standard  
 535 Model. It provides a solution to the hierarchy problem, leads to gauge coupling  
 536 unification, and provides a dark matter candidate consistent with galactic rotation  
 537 curves. As noted in this chapter, due to the LSPs in the final state, most SUSY  
 538 searches require a significant amount of missing transverse energy in combination  
 539 with jets of high transverse momentum. However, there is some opportunity to do  
 540 better than this, especially in final states where one has two weakly-interacting LSPs  
 541 on opposite sides of some potentially complicated decay tree. We will see how this is  
 542 done in Ch.??.



*The Large Hadron Collider*

545 The Large Hadron Collider (LHC) produces high-energy protons which are collided  
 546 at the center of multiple large experiments at CERN on the outskirts of Geneva,  
 547 Switzerland [85]. The LHC produces the highest energy collisions in the world,  
 548 with design center-of-mass energy of  $\sqrt{s} = 14$  TeV, which allows the experiments  
 549 to investigate physics far beyond the reach of previous colliders. This chapter will  
 550 summarize the basics of accelerator physics, especially with regards to discovering  
 551 physics beyond the Standard Model. We will describe the CERN accelerator complex  
 552 and the LHC.

553 **4.1 Basics of Accelerator Physics**

554 This section follows closely the presentation of [86].

Simple particle accelerators simply rely on the acceleration of charged particles in a static electric field. Given a field of strength  $E$ , charge  $q$ , and mass  $m$ , this is simply

$$a = \frac{qE}{m}. \quad (4.1)$$

555 For a given particle with a given mass and charge, this is limited by the static electric  
 556 field which can be produced, which in turn is limited by electrical breakdown at high  
 557 voltages.

558 There are two complementary solutions to this issue. First, we use the *radio*  
 559 *frequency acceleration* technique. We call the devices used for this *RF cavities*. The

560    cavities produce a time-varied electric field, which oscillate such that the charged  
 561    particles passing through it are accelerated towards the design energy of the RF  
 562    cavity. This oscillation also induces the particles into *bunches*, since particles which  
 563    are slightly off in energy from that induced by the RF cavity are accelerated towards  
 564    the design energy.

Second, one bends the particles in a magnetic field, which allows them to pass through the same RF cavity over and over. This second process is often limited by *synchrotron radiation*, which describes the radiation produced when a charged particle is accelerated. The power radiated is

$$P \sim \frac{1}{r^2} \left( E/m \right)^4 \quad (4.2)$$

565    where  $r$  is the radius of curvature and  $E, m$  is the energy (mass) of the charged  
 566    particle. Given an energy which can be produced by a given set of RF cavities (which  
 567    is *not* limited by the mass of the particle), one then has two options to increase the  
 568    actual collision energy : increase the radius of curvature or use a heavier particle.  
 569    Practically speaking, the easiest options for particles in a collider are protons and  
 570    electrons, since they are (obviously) copious in nature and do not decay<sup>1</sup>. Given the  
 571    dependence on mass, we can see why protons are used to reach the highest energies.  
 572    The tradeoff for this is that protons are not point particles, and we thus we don't  
 573    know the exact incoming four-vectors of the protons, as discussed in Ch.2.

The particle *beam* refers to the bunches all together. An important property of a beam of a particular energy  $E$ , moving in uniform magnetic field  $B$ , containing particles of momentum  $p$  is the *beam rigidity* :

$$R \equiv rB = p/c. \quad (4.3)$$

574    The linear relation between  $r$  and  $p$ , or alternatively  $B$  and  $p$  have important  
 575    consequences for LHC physics. For hadron colliders, this is the limiting factor on

---

<sup>1</sup>Muon colliders are a really cool option at high energies, since the relativistic  $\gamma$  factor gives them a relatively long lifetime in the lab frame.

576 going to higher energy scales; one needs a proportionally larger magnetic field to  
577 keep the beam accelerating in a circle.

578 Besides the rigidity of the beam, the most important quantities to characterize  
579 a beam are known as the (normalized) *emittance*  $\epsilon_N$  and the *betatron function*  $\beta$ .  
580 These quantities determine the transverse size  $\sigma$  of a relativistic beam  $v \gtrsim c$  beam :  
581  $\sigma^2 = \beta^* \epsilon_N / \gamma_{\text{rel}}$ , where  $\beta^*$  is the value of the betatron function at the collision point  
582 and  $\gamma_{\text{rel}}$  is the Lorentz factor.

These quantities determine the *instantaneous luminosity*  $L$  of a collider, which combined with the cross-section  $\sigma$  of a particular physics process, give the rate of this physics process :

$$R = L\sigma. \quad (4.4)$$

The instantaneous luminosity  $L$  is given by :

$$L = \frac{f_{\text{rev}} N_b^2 F}{4\pi\sigma^2} = \frac{f_{\text{rev}} n N_b^2 \gamma_{\text{rel}} F}{4\pi\beta^* \epsilon_N}. \quad (4.5)$$

583 Here we have introduced the frequency of revolutions  $f_{\text{rev}}$ , the number of bunches  $n$ ,  
584 the number of protons per bunch  $N_b^2$ , and a geometric factor  $F$  related to the crossing  
585 angle of the beams.

The *integrated luminosity*  $\int L$  gives the total number of a particular physics process  $P$ , with cross-section  $\sigma_P$ .

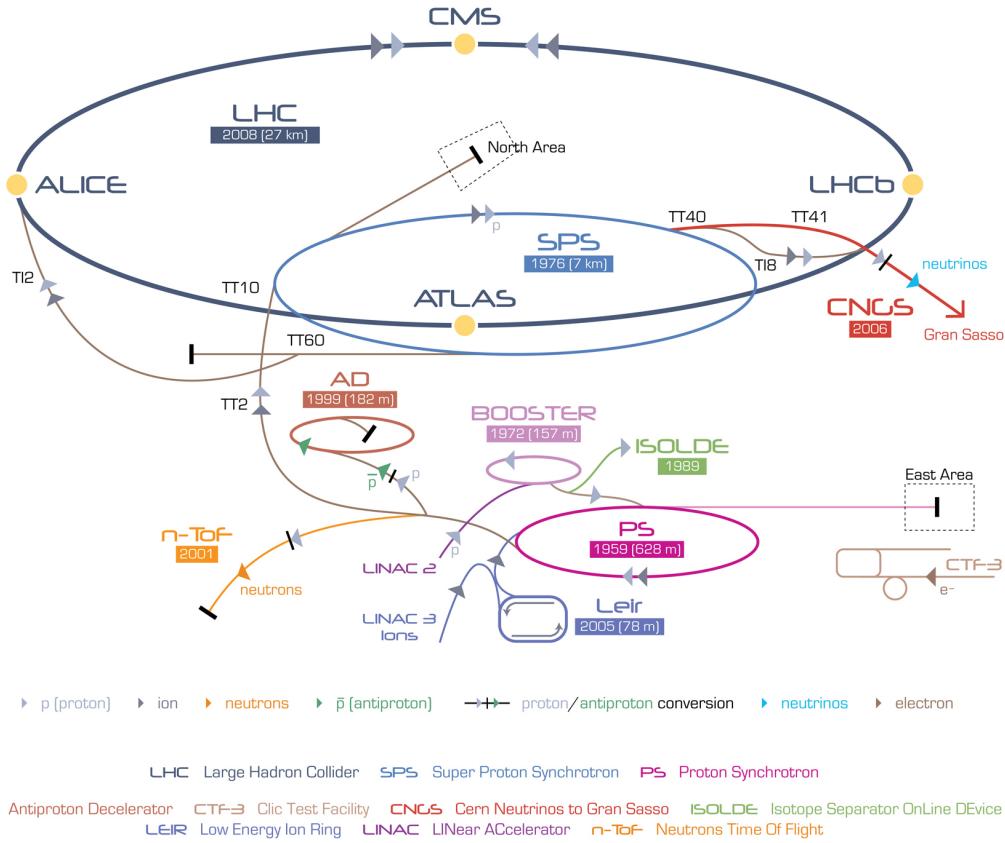
$$N_P = \sigma_P \int L. \quad (4.6)$$

586 Due to this simple relation, one can also quantify the “amount of data delivered” by  
587 a collider simply by  $\int L$ .

## 588 4.2 Accelerator Complex

589 The Large Hadron Collider is the last accelerator in a chain of accelerators which  
590 together form the CERN accelerator complex, which can be seen in 4.1. The protons

Figure 4.1: The CERN accelerator complex.



begin their journey to annihilation in a hydrogen source, where they are subsequently ionized. The first acceleration occurs in the Linac 2, a linear accelerator composed of RF cavities. The protons leave the Linac 2 at an energy of 50 MeV and enter the Proton Synchrotron Booster (PSB). The PSB contains four superimposed rings, which accelerate the protons to 1.4 GeV. The protons are then injected into the Proton Synchrotron (PS). This synchrotron increases the energy up to 25 GeV. After leaving the PS, the protons enter the Super Proton Synchrotron (SPS). This is the last step before entering the LHC ring, and the protons are accelerated to 450 GeV. From the SPS, the protons are injected into the beam pipes of the LHC. The process to fill the LHC rings with proton bunches from start to finish typically takes about four minutes.

602 **4.3 Large Hadron Collider**

The Large Hadron Collider is the final step in the CERN accelerator complex, and produces the collisions analyzed in this thesis. From the point of view of experimentalists on the general-purpose ATLAS and CMS experiments, the main goal of the LHC is to deliver collisions at the highest possible energy, with the highest possible instantaneous luminosity. The LHC was installed in the existing 27 km tunnel used by the Large Electron Positron (LEP) collider [87]. This allowed the existing accelerator complex at CERN, described in the previous section, to be used as the injection system to prepare the protons up to 450 GeV. Many aspects of the LHC design were decided by this very fact, and specified the options allowed to increase the energy or luminosity. In particular, the radius of the tunnel was already specified; from Eq.4.3, this implies the momentum (or energy) of the beam is entirely determined by the magnetic field. Given the 27 km circumference of the LEP tunnel, one can calculate the required magnetic field to reach the 7 TeV per proton design energy of the LHC :

$$r = C/2\pi = 4.3 \text{ km} \quad (4.7)$$

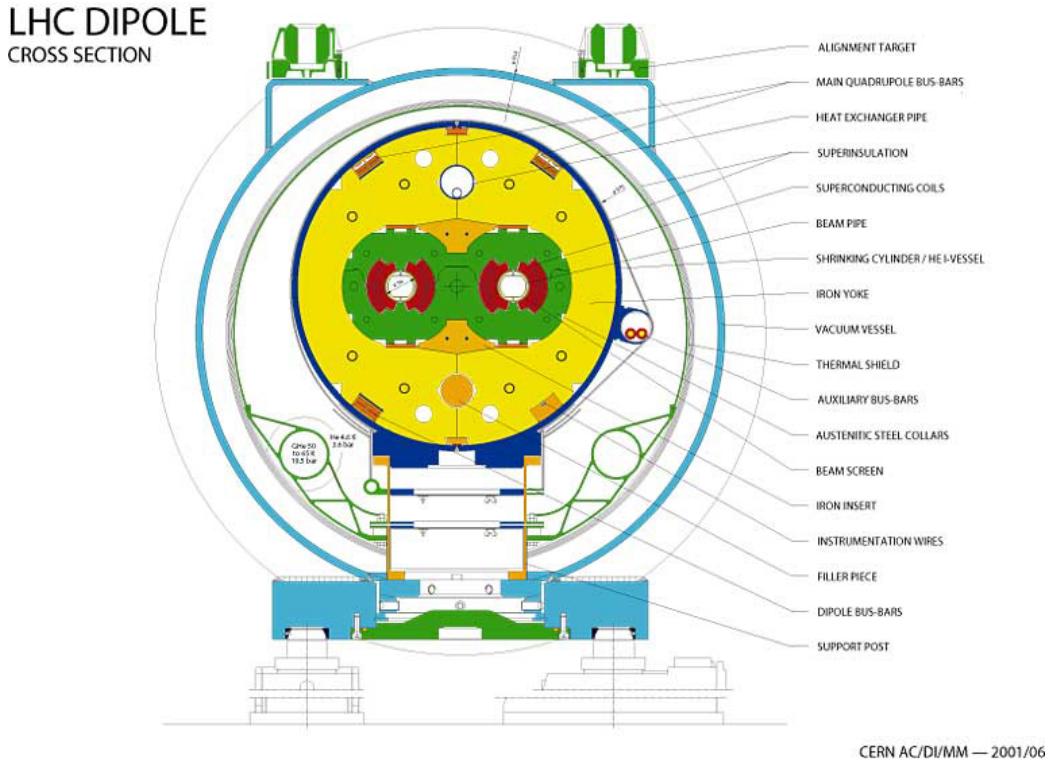
$$\rightarrow B = \frac{p}{rc} = 5 \text{ T} \quad (4.8)$$

603 In fact, the LHC consists of 8 528 m straight portions consisting of RF cavities, used  
604 to accelerate the particles, and 8 circular portions which bend the protons around the  
605 LHC ring. These circular portions actually have a slightly smaller radius of curvature  
606  $r = 2804 \text{ m}$ , and we require  $B = 8.33 \text{ T}$ . To produce this large field, we need to use  
607 superconducting magnets, as discussed in the next section.

608 **Magnets**

609 There are many magnets used by the LHC machine, but the most important are the  
610 1232 dipole magnets; a schematic is shown in Fig.4.2 and a photograph is shown in

Figure 4.2: Schematic of an LHC dipole magnet.



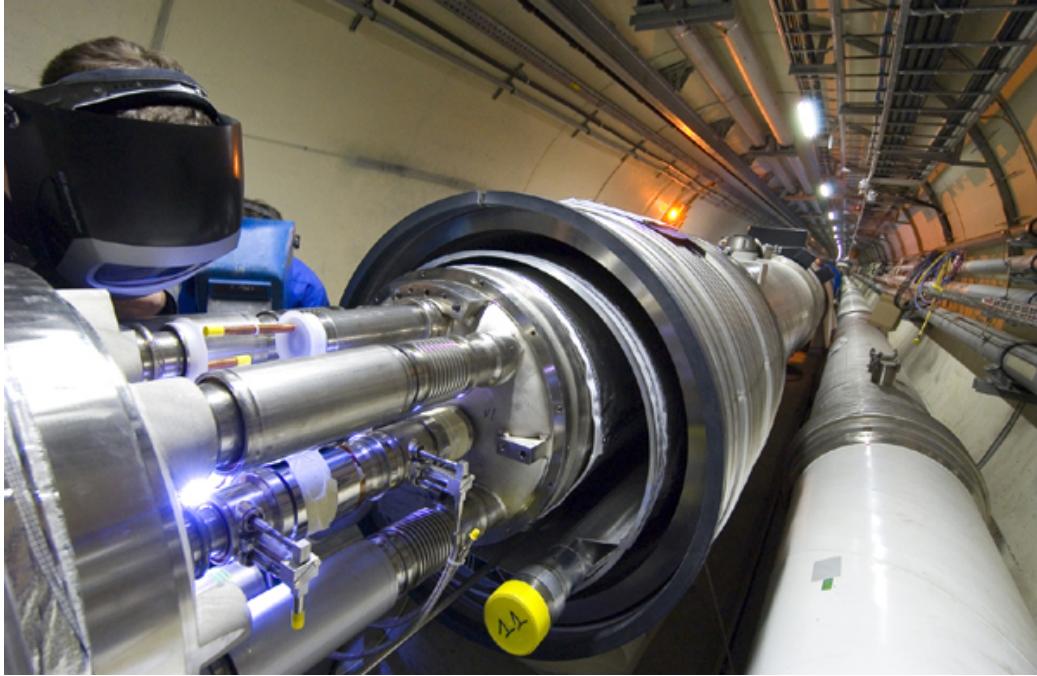
611 Fig.4.3.

612 The magnets are made of Niobium and Titanium. The maximum field strength is  
 613 10 T when cooled to 1.9 Kelvin. The magnets are cooled by superfluid helium, which  
 614 is supplied by a large cryogenic system. Due to heating between the eight helium  
 615 refrigerators and the beampipe, the helium is cooled in the refrigerators to 1.8 K.

616 A failure in the cooling system can cause what is known as a *quench*. If the  
 617 temperature goes above the critical superconducting temperature, the metal loses its  
 618 superconducting properties, which leads to a large resistance in the metal. This leads  
 619 to rapid temperature increases, and can cause extensive damages if not controlled.

620 The dipole magnets are 16.5 meters long with a diameter of 0.57 meters. There  
 621 are two individual beam pipes inside each magnet, which allows the dipoles to house  
 622 the beams travelling in both directions around the LHC ring. They curve slightly,  
 623 at an angle of 5.1 mrad, which carefully matches the curvature of the ring. The

Figure 4.3: Photograph of a technician connecting an LHC dipole magnet.



624 beampipes inside of the magnets are held in high vacuum, to avoid stray particles  
625 interacting with the beam.

## 626 **4.4 Dataset Delivered by the LHC**

627 In this thesis, we analyze the data delivered by the LHC to ATLAS in the 2015 and  
628 2016 datasets. The beam parameters relevant to this dataset are available in Table  
629 [4.1](#).

630 The peak instantaneous luminosity delivered in 2015 (2016) was  $L =$   
631  $5.2(11) \text{ cm}^{-2}\text{s}^{-1} \times 10^{33}$ . One can note that the instantaneous luminosity delivered in  
632 the 2016 dataset exceeds the design luminosity of the LHC. The total integrated  
633 luminosity delivered was  $13.3 \text{ fb}^{-1}$ . In Figure [4.4](#), we display the integrated luminosity  
634 as a function of day for 2015 and 2016.

Figure 4.4: Integrated Luminosity delivered by the LHC and collected by ATLAS in the 2015 and 2016 datasets.

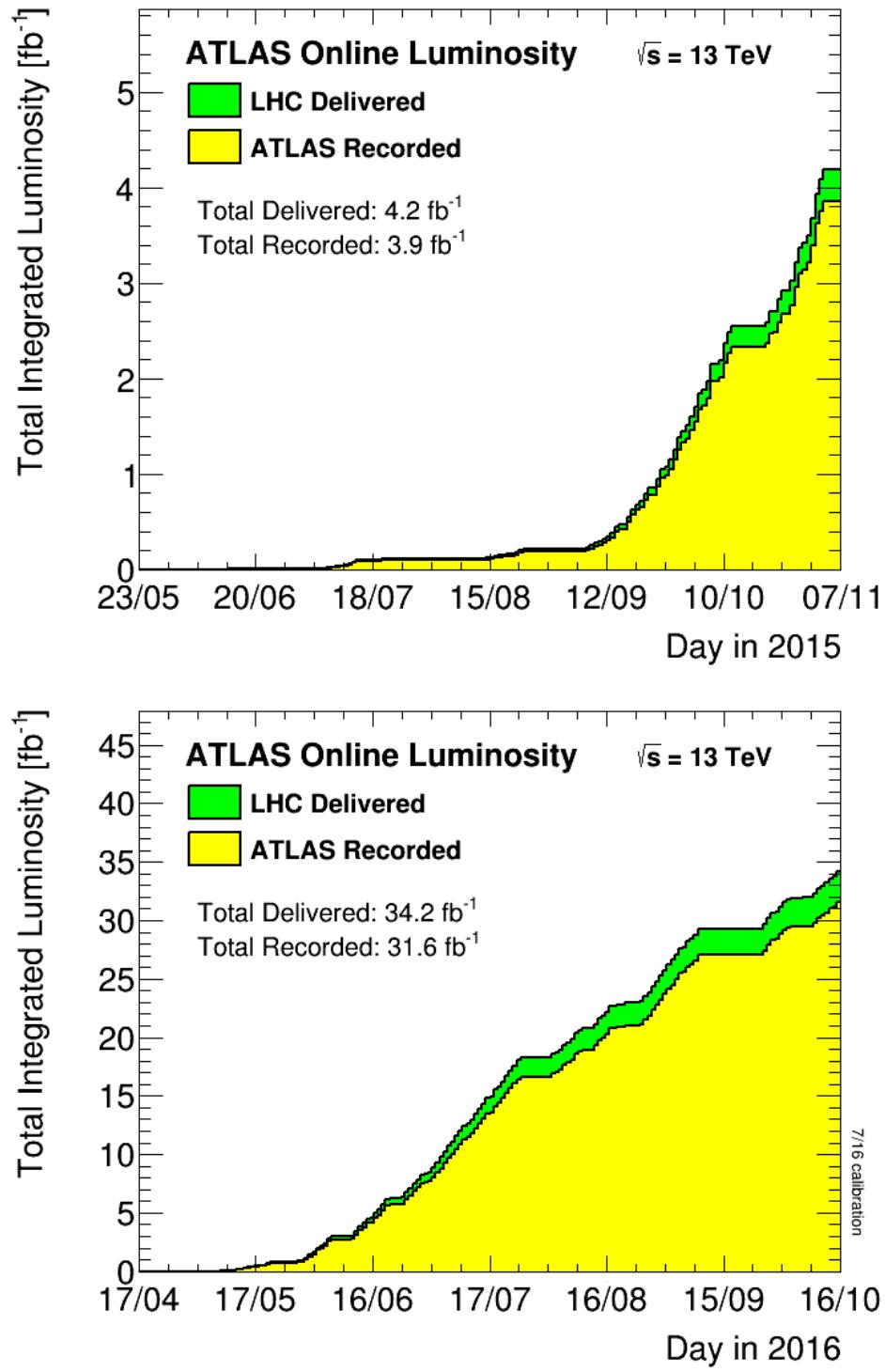
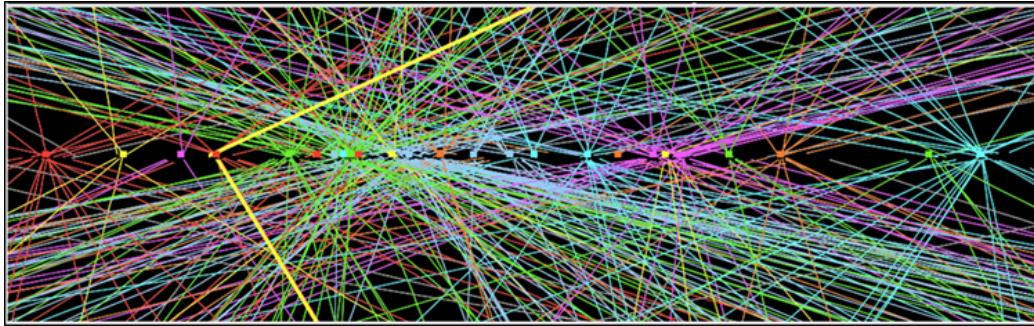


Table 4.1: Beam parameters of the Large Hadron Collider.

Parameter	Injection	Extraction
Energy (GeV)	450	7000
Rigidity (T-m)	3.8	23353
Bunch spacing (ns)	25	25
Design Luminosity ( $\text{cm}^{-2}\text{s}^{-1} \times 10^{34}$ )	-	1.0
Bunches per proton beam	2808	2808
Protons per bunch	1.15 e11	1.15 e11
Beam lifetime (hr)	-	10
Normalized Emittance $\epsilon_N$ (mm $\mu\text{rad}$ )	3.3	3.75
Betatron function at collision point $\beta^*$ (cm)	-	55

Figure 4.5: Simulated event with many pileup vertices.



## 635 Pileup

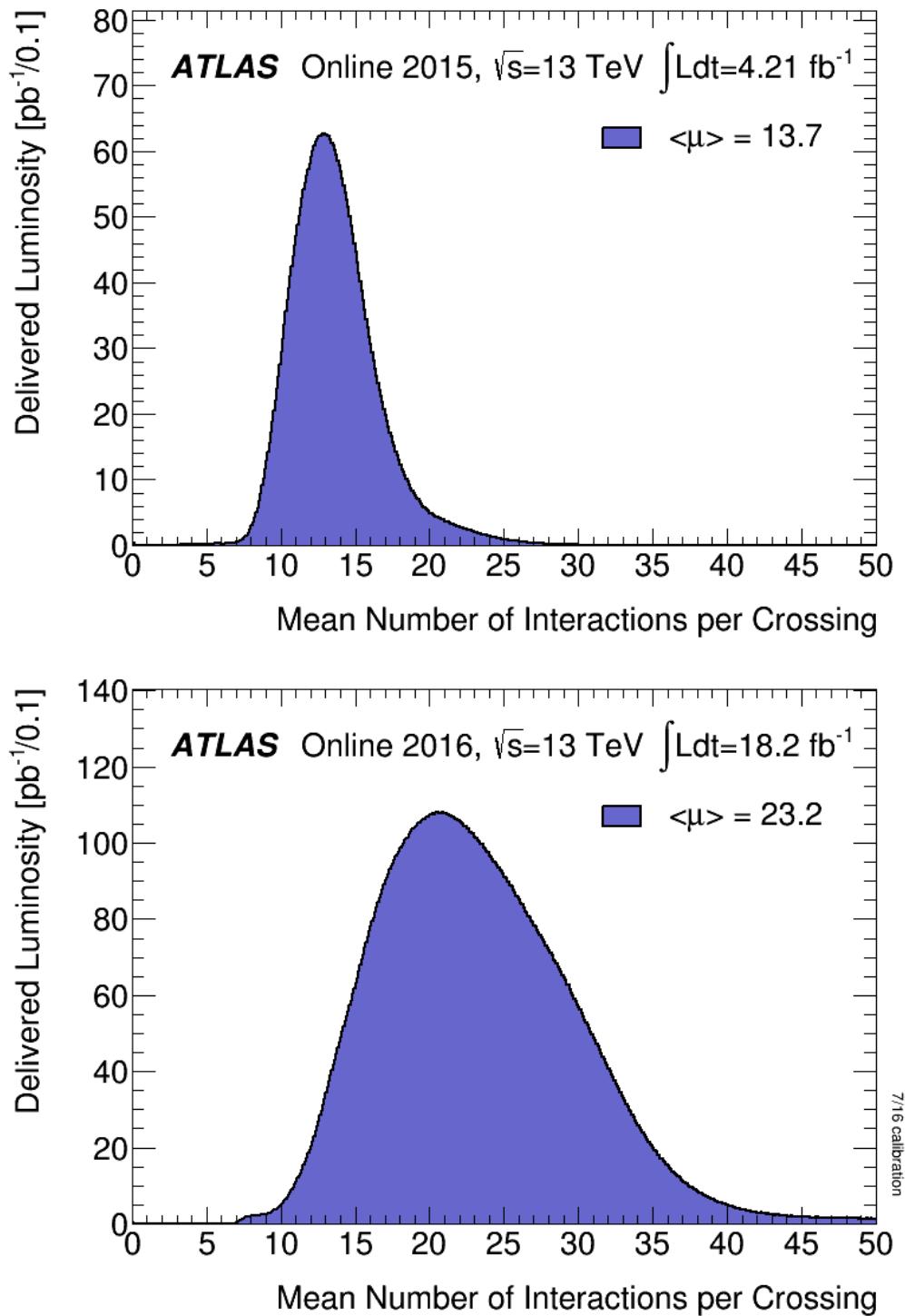
636 *Pileup* is the term for the additional proton-proton interactions which occur during  
 637 each bunch crossing of the LHC. At the beginning of the LHC physics program, there  
 638 had not been a collider which averaged more than a single interaction per bunch  
 639 crossing. In the LHC, each bunch crossing (or *event*) generally contains multiple  
 640 proton-proton interactions. An simulated event with many *vertices* can be seen in  
 641 Fig.4.5. The so-called *primary vertex* (or *hard scatter vertex*) refers to the vertex  
 642 which has the highest  $\Sigma p_T^2$ ; this summation occurs over the *tracks* in the detector,  
 643 which we will describe later[**ATL-INDET-PUB-2009-001**]. We then distinguish  
 644 between *in-time* pileup and *out-of-time* pileup. In-time pileup refers to the additional  
 645 proton-proton interactions which occur in the event. Out-of-time pileup refers to  
 646 effects related to proton-proton interactions previous bunch crossings.

647        We quantify in-time pileup by the number of “primary”<sup>2</sup> vertices in a particular  
648    event. To quantify the out-of-time pileup, we use the average number of interactions  
649    per bunch crossing  $\langle \mu \rangle$  over some human-scale time. In Figure 4.6, we show the  
650    distribution of  $\mu$  for the dataset used in this thesis.

---

<sup>2</sup>The primary vertex is as defined above, but we unfortunately use the same name here.

Figure 4.6: Mean number of interactions per bunch crossing in the 2015 and 2016 datasets.





*The ATLAS detector*

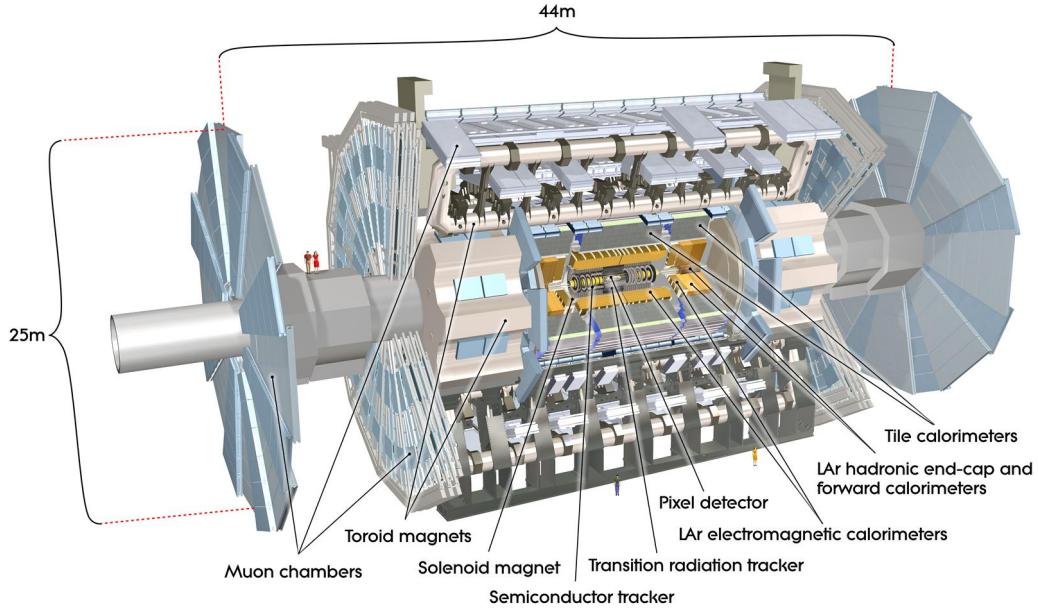
653 The dataset analyzed in this thesis was taken by the ATLAS detector [88], which is  
 654 located at the “Point 1” cavern of the LHC beampipe, just across the street from  
 655 the main CERN campus. The much-maligned acronym stands for *A Toroidal LHC*  
 656 *ApparatuS*. ATLAS is a massive cylindrical detector, with a radius of 12.5 m and a  
 657 length of 44 m, with nearly hermitic coverage around the collision point. It consists  
 658 of multiple subdetectors; each plays a role in ATLAS’s ultimate purpose of measuring  
 659 the energy, momentum, and type of the particles produced in collisions delivered by  
 660 the LHC. These subdetectors are immersed in a hybrid solenoid-toroid magnet system  
 661 whichs forces charged particles to curve, which allows for precise measurements of  
 662 their momenta. These magnetic fields are maximized in the central solenoid magnet,  
 663 which contains a magnetic field of 2 T. A schematic of the detector can be seen in  
 664 [5.1](#).

665 The *inner detector* (ID) lies closest to the collision point, and contains three  
 666 separate subdetectors. It provides pseudorapidity<sup>1</sup>coverage of  $|\eta| < 2.5$  for charged  
 667 particles to interact with the tracking material. The tracks reconstructed from the  
 668 inner detector hits are used to reconstruct the primary vertices, as noted in Ch.??,

---

<sup>1</sup>ATLAS uses a right-handed Cartesian coordinate system; the origin is defined by the nominal beam interaction point. The positive- $z$  direction is defined by the incoming beam travelling counterclockwise around the LHC. The positive- $x$  direction points towards the center of the LHC ring from the origin, and the positive- $y$  direction points upwards towards the sky. For particles of transverse (in the  $x - y$  plane) momentum  $p_T = \sqrt{p_x^2 + p_y^2}$  and energy  $E$ , it is generally most convenient fully describe this particle’s kinematics as measured by the detector in the  $(p_T, \phi, \eta, E)$  basis. The angle  $\phi = \arctan(p_y/p_x)$  is the standard azimuthal angle, and  $\eta = \ln \tan(\theta/2)$  is known as the pseudorapidity, and defined based on the standard polar angle  $\theta = \arccos(p_z/p_T)$ . For locations of i.e. detector elements, both  $(r, \phi, \eta)$  and  $(z, \phi, \eta)$  can be useful.

Figure 5.1: The ATLAS detector

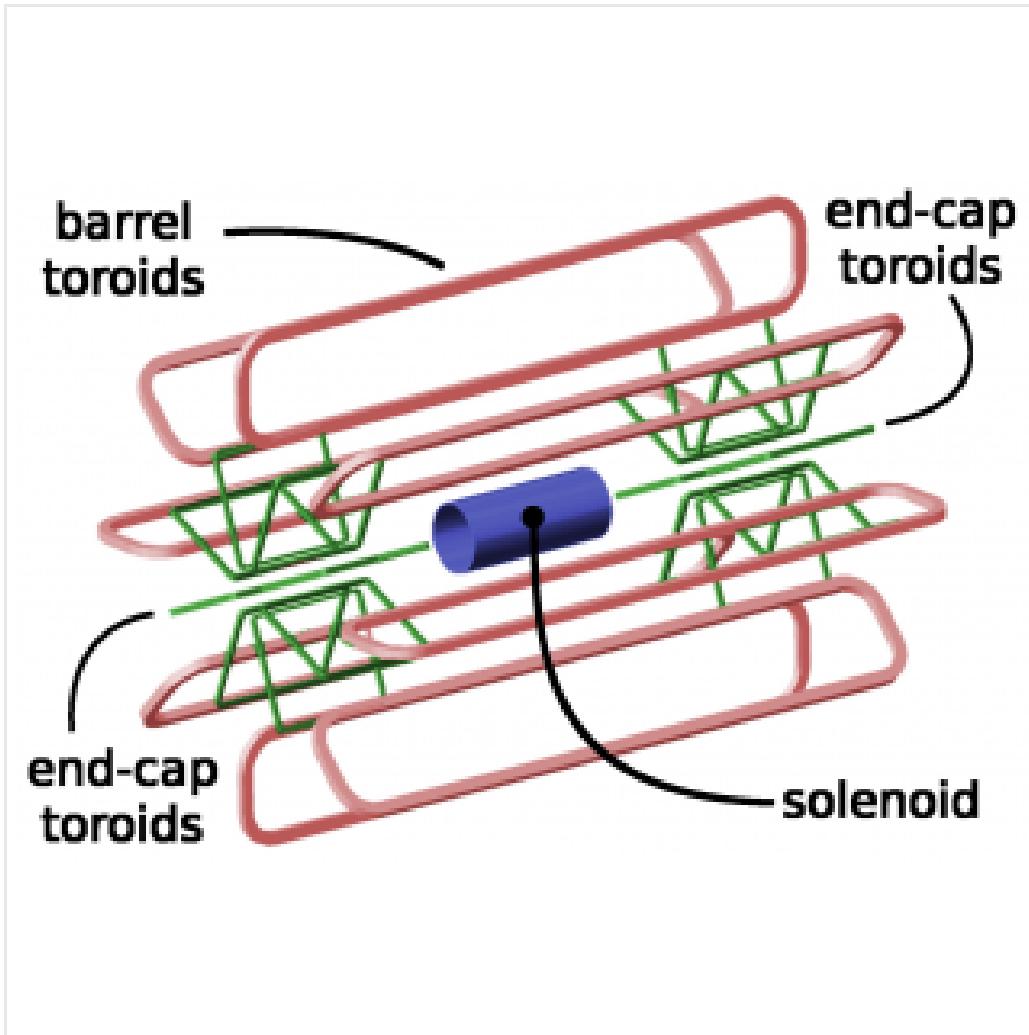


and to determine the momenta of charged particles. The ATLAS *calorimeter* consists of two subdetectors, known as the *electromagnetic* and *hadronic* calorimeters. These detectors stop particles in their detector material, and measure the energy deposition inside, which measures the energy of the particles deposited. The calorimeters provide coverage out to pseudorapidity of  $|\eta| < 4.9$ . The muon spectrometer is aptly named; it is specifically used for muons, which are the only particles which generally reach the outer portions of the detector. In this region, we have the large tracking systems of the muon spectrometer, which provide precise measurements of muon momenta. The muon spectrometer has pseudorapidity coverage of  $|\eta| < 2.7$ .

## 5.1 Magnets

ATLAS contains multiple magnetic systems; primarily, we are concerned with the solenoid, used by the inner detector, and the toroids located outside of the ATLAS calorimeter. A schematic is shown in Fig.5.2. These magnetic fields are used to bend

Figure 5.2: The ATLAS magnet system



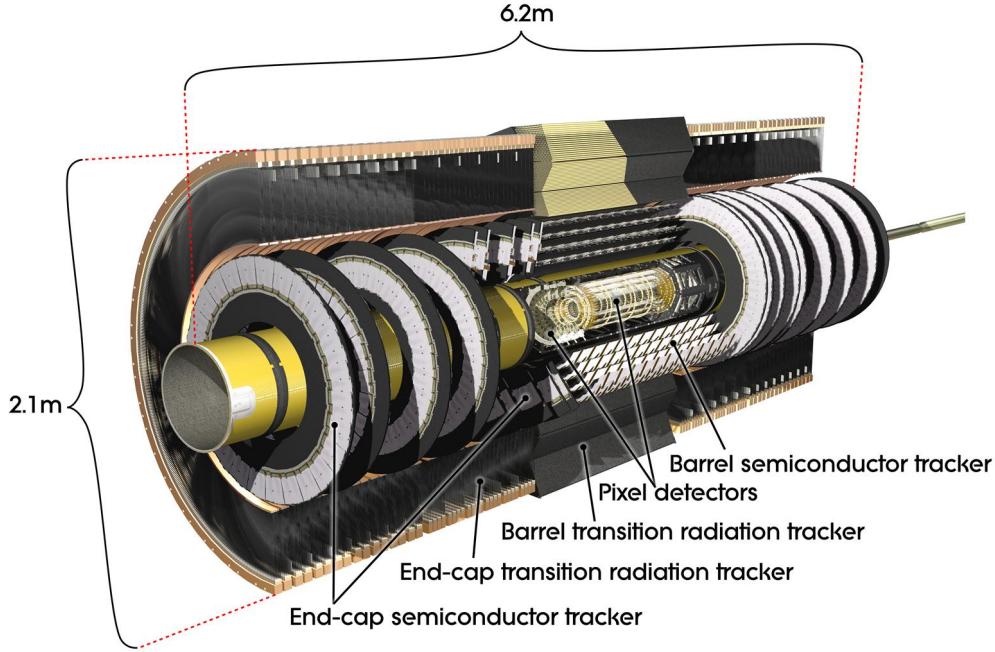
682 charged particles under the Lorentz force, which subsequently allows one to measure  
683 their momentum.

684 The ATLAS central solenoid is a 2.3 m diameter, 5.3 m long solenoid at the center  
685 of the ATLAS detector. It produces a uniform magnetic field of 2 T; this strong field  
686 is necessary to accurately measure the charged particles in this field. An important  
687 design constraint for the central solenoid was the decision to place it in between the  
688 inner detector and the calorimeters. To avoid excessive impacts on measurements in  
689 the calorimetry, the central solenoid must be as transparent as possible<sup>2</sup>.

---

<sup>2</sup>This is also one of the biggest functional differences between ATLAS and CMS; in CMS, the

Figure 5.3: The ATLAS inner detector



690     The toroid system consists of eight air-core superconducting barrel loops; these  
691    give ATLAS its distinctive shape. There are also two endcap air-core magnets. These  
692    produce a magnetic field in a region of approximately 26 m in length and 10 m of  
693    radius. The magnetic field in this region is non-uniform, due to the prohibitive costs  
694    of a solenoid magnet of that size.

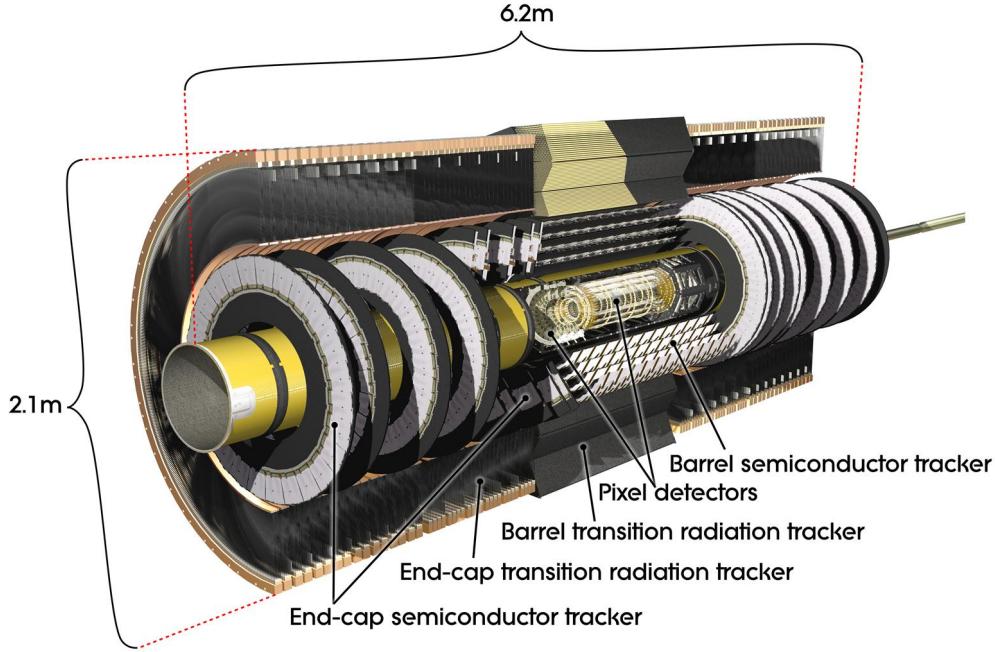
## 695    **5.2 Inner Detector**

696    The ATLAS inner detector consists of three separate tracking detectors, which are  
697    known as, in order of increasing distance from the interaction point, the Pixel  
698    Detector, Semiconductor Tracker (SCT), and the Transition Radiation Tracker  
699    (TRT). When charged particles pass through these tracking layers, they produce  
700    *hits*, which using the known 2 T magnetic field, allows the reconstruction of *tracks*.  
701    Tracks are used as inputs for reconstruction of many higher-level physics objects,

---

solenoid is outside of the calorimeters.

Figure 5.4: The ATLAS pixel detector



such as electrons, muons, photons, and  $E_T^{\text{miss}}$ . Accurate track reconstruction is thus crucial for precise measurements of charged particles.

## Pixel Detector

The ATLAS pixel detector consists four layers of silicon “pixels”. This refers to the segmentation of the active medium into the pixels; compare to the succeeding silicon detectors, which will use silicon “strips”. This provides precise 3D hit locations. The layers are known as the “Insertable”<sup>3</sup>B-Layer (IBL), the B-Layer (or Layer-0), Layer-1, and Layer-2, in order of increasing distance from the interaction point. These layers are very close to the interaction point, and therefore experience a large amount of radiation.

Layer-1, Layer-2, and Layer-3 were installed with the initial construction of ATLAS. They contain front-end integrated electronics (FEI3s) bump-bonded to 1744

---

<sup>3</sup>Very often, the IBL is mistakenly called the Inner B-Layer, which would have been a much more sensible name.

714 silicon modules; each module is  $250 \mu\text{m}$  in thickness and contains 47232 pixels. These  
715 pixels have planar sizes of  $50 \times 400 \mu\text{m}^2$  or  $50 \times 600 \mu\text{m}^2$ , to provide highly accurate  
716 location information. The FEI3s are mounted on long rectangular structures known  
717 as staves, which encircle the beam pipe. A small tilt to each stave allows full coverage  
718 in  $\phi$  even with readout systems which are installed. These layers are at radia of 50.5  
719 mm, 88.5 mm, and 122.5 mm from the interaction point.

720 The IBL was added to ATLAS after Run1 in 2012 at a radius of 33 mm from the  
721 interaction point. The entire pixel detector was removed from the center of ATLAS  
722 to allow an additional pixel layer to be installed. The IBL was required to preserve  
723 the integrity of the pixel detector as radiation damage leads to inoperative pixels in  
724 the other layers. The IBL consists of 448 FEI4 chips, arranged onto 14 staves. Each  
725 FEI4 has 26880 pixels, of planar size  $50 \times 250 \mu\text{m}$ . This smaller granularity was  
726 required due to the smaller distance to the interaction point.

727 In total, a charged particle passing through the inner detector would expect to  
728 leave four hits in the pixel detector.

## 729 Semiconductor Tracker

730 The SCT is directly beyond Layer-2 of the pixel detector. This is a silicon strip  
731 detector, which do not provide the full 3D information of the pixel detector. The  
732 dual-sensors of the SCT contain  $2 \times 768$  individual strips; each strip has area  $6.4 \text{ cm}^2$ .  
733 The SCT dual-sensor is then double-layered, at a relative angle of 40 mrad;  
734 together these layers provide the necessary 3D information for track reconstruction.  
735 There are four of these double-layers, at radia of 284 mm, 355 mm, 427 mm, and 498  
736 mm. These double-layers provide hits comparable to those of the pixel detector, and  
737 we have four additional hits to reconstruct tracks for each charged particle.

Figure 5.5: A ring of the Semiconductor Tracker

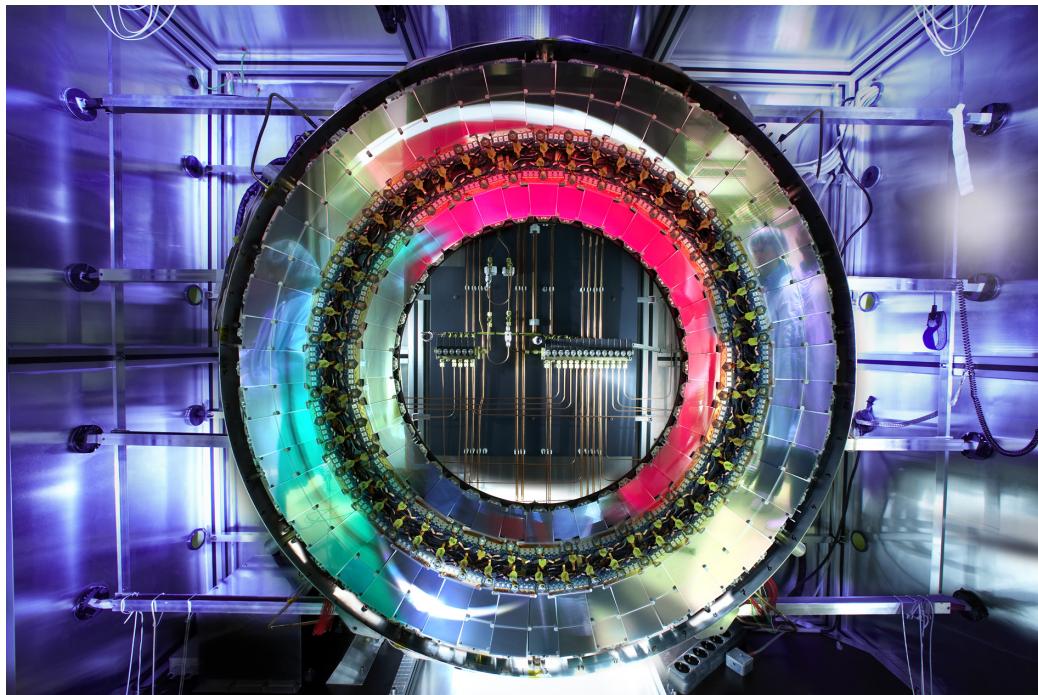
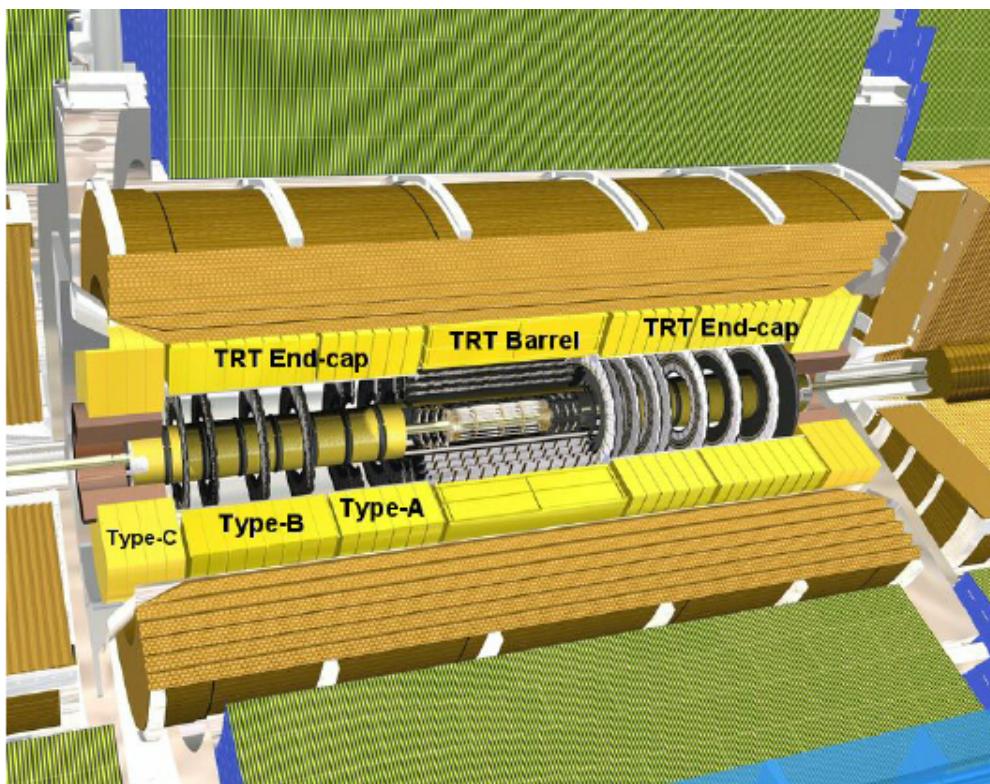


Figure 5.6: A schematic of the Transition Radiation Tracker



738 **Transition Radiation Tracker**

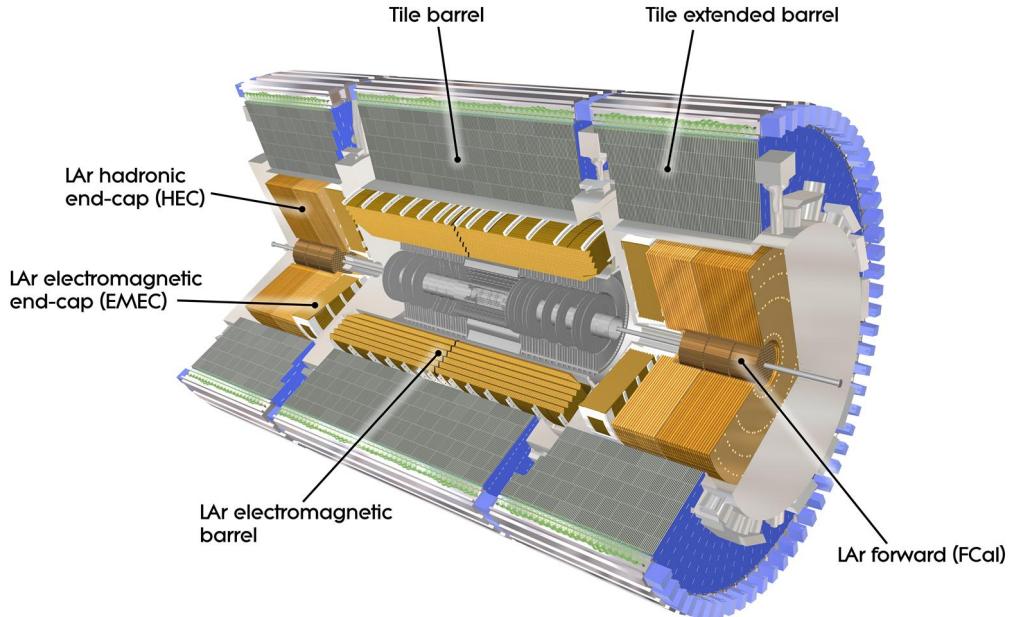
739 The Transition Radiation Tracker is the next detector radially outward from the SCT.  
740 It contains straw drift tubes; these contain a tungsten gold-plated wire of  $32 \mu\text{m}$   
741 diameter held under high voltage (-1530 V) with the edge of the Kapton-aluminum  
742 tube. They are filled with a gas mixture of primarily xenon that is ionized when  
743 a charged particle passes through the tube. The ions are collected by the “drift”  
744 due to the voltage inside the tubes, which is read out by the electronics. This gives  
745 so-called “continuous tracking” throughout the tube, due to the large number of ions  
746 produced.

747 The TRT is so-named due to the *transition radiation* (TR) it induces. Due to  
748 the dielectric difference between the gas and tubes, TR is induced. This is important  
749 for distinguishing electrons from their predominant background of minimum ionizing  
750 particles. Generally, electrons have a much larger Lorentz factor than minimum  
751 ionizing particles, which leads to additional TR. This can be used as an additional  
752 handle for electron reconstruction.

753 **5.3 Calorimetry**

754 The calorimetry of the ATLAS detector also includes multiple subdetectors; these sub-  
755 detectors allow precise measurements of the electrons, photons, and hadrons produced  
756 by the ATLAS detector. Generically, calorimeters work by stopping particles in their  
757 material, and measuring the energy deposition. This energy is deposited as a cascade  
758 particles induce from interactions with the detector material known *showers*. ATLAS  
759 uses *sampling* calorimeters; these alternate a dense absorbing material, which induces  
760 showers, with an active layer which measures energy depositions by the induced  
761 showers. Since some energy is deposited into the absorption layers as well, the energy  
762 depositions must be properly calibrated for the detector.

Figure 5.7: The ATLAS calorimeter

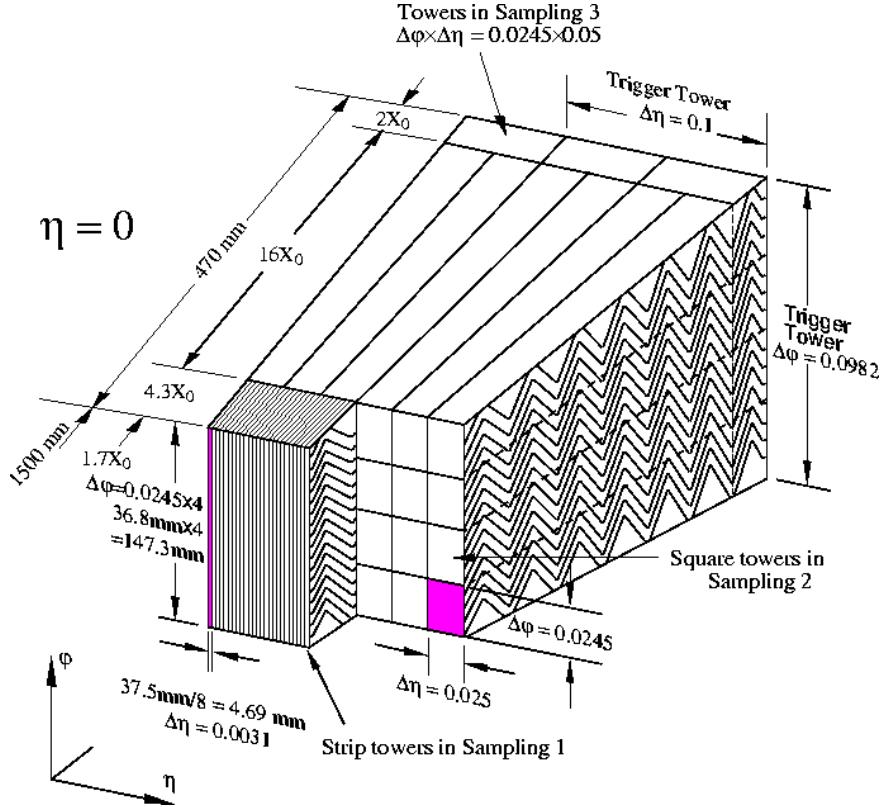


763 Electromagnetic objects (electrons and photons) and hadrons have much different  
764 interaction properties, and thus we need different calorimeters to accurately measure  
765 these different classes of objects; we can speak of the *electromagnetic* and *hadronic*  
766 calorimeters. ATLAS contains four separate calorimeters : the liquid argon (LAr)  
767 electromagnetic barrel calorimeter, the Tile barrel hadronic calorimeter, the LAr  
768 endcap electromagnetic calorimeter, the LAr endcap hadronic calorimeter, and the  
769 LAr Forward Calorimeter (FCal). Combined, these provide full coverage in  $\phi$  up to  
770  $|\eta| < 4.9$ , and can be seen in Fig.5.7.

## 771 **Electromagnetic Calorimeters**

772 The electromagnetic calorimeters of the ATLAS detector consist of the barrel and  
773 endcap LAr calorimeters. These are arranged into an ingenious “accordion” shape,  
774 shown in 5.8, which allows full coverage in  $\phi$  and exceptional coverage in  $\eta$  while  
775 still allowing support structures for detector operation. The accordion is made of

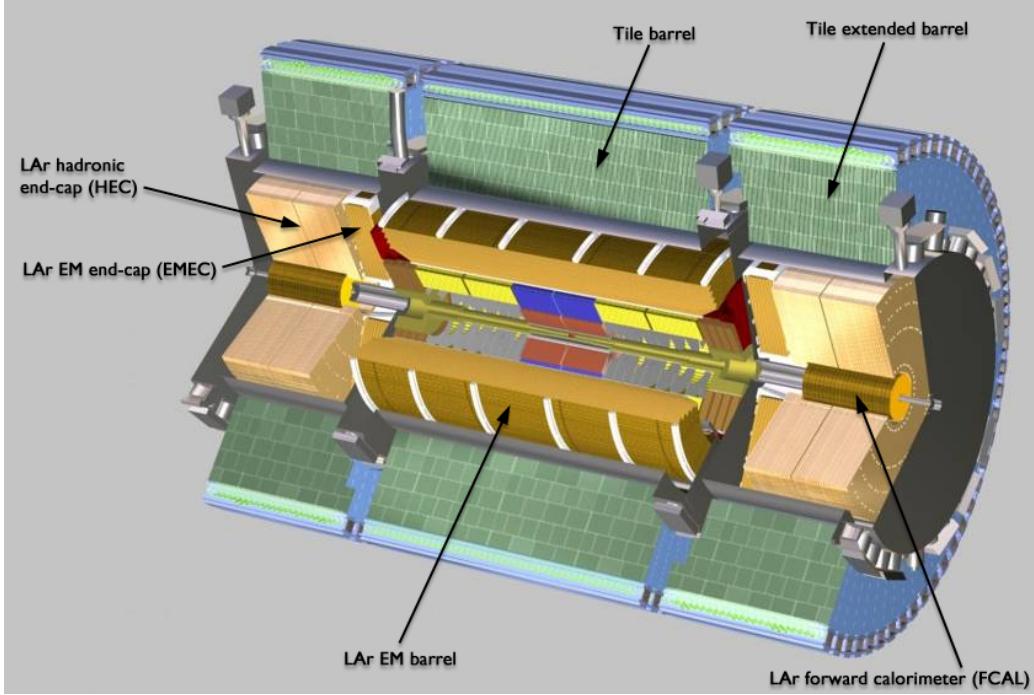
Figure 5.8: A schematic of a subsection of the barrel LAr electromagnetic calorimeter



776 layers with liquid argon (active detection material) and lead (absorber) to induce  
 777 electromagnetic showers. The LAr EM calorimeters are each more than 20 radiation  
 778 lengths deep, which provides the high stopping power necessary to properly measure  
 779 the electromagnetic showers.

780 The barrel component of the LAr EM calorimeter extends from the center of the  
 781 detector out to  $|\eta| < 1.475$ . The calorimeter has a presampler, which measures the  
 782 energy of any EM shower induced before the calorimeter. This has segmentation of  
 783  $\Delta\eta = 0.025, \Delta\phi = .01$ . There are three “standard” layers in the barrel, which have  
 784 decreasing segmentation into calorimeter *cells* as one travels radially outward from  
 785 the interaction point. The first layer has segmentation of  $\Delta\eta = 0.003, \Delta\phi = .1$ , and  
 786 is quite thin relative to the other layers at only 4 radiation lengths deep. It provides  
 787 precise  $\eta$  and  $\phi$  measurements for incoming EM objects. The second layer is the  
 788 deepest at 16 radiation lengths, with a segmentation of  $\Delta\eta = 0.025, \Delta\phi = 0.025$ . It

Figure 5.9: A schematic of Tile hadronic calorimeter



is primarily responsible for stopping the incoming EM particles, which dictates its large relative thickness, and measures most of the energy of the incoming particles. The third layer is only 2 radiation lengths deep, with a rough segmentation of  $\Delta\eta = 0.05$ ,  $\Delta\phi = .025$ . The deposition in this layer is primarily used to distinguish hadrons interacting electromagnetically and entering the hadronic calorimeter from the strictly EM objects which are stopped in the second layer.

The barrel EM calorimeter has a similar overall structure, but extends from  $1.4 < |\eta| < 3.2$ . The segmentation in  $\eta$  is better in the endcap than the barrel; the  $\phi$  segmentation is the same. In total, the EM calorimeters contain about 190000 individual calorimeter cells.

## Hadronic Calorimeters

The hadronic calorimetry of ATLAS sits directly outside the EM calorimetry. It contains three subdetectors : the barrel Tile calorimeter, the endcap LAr calorimeter,

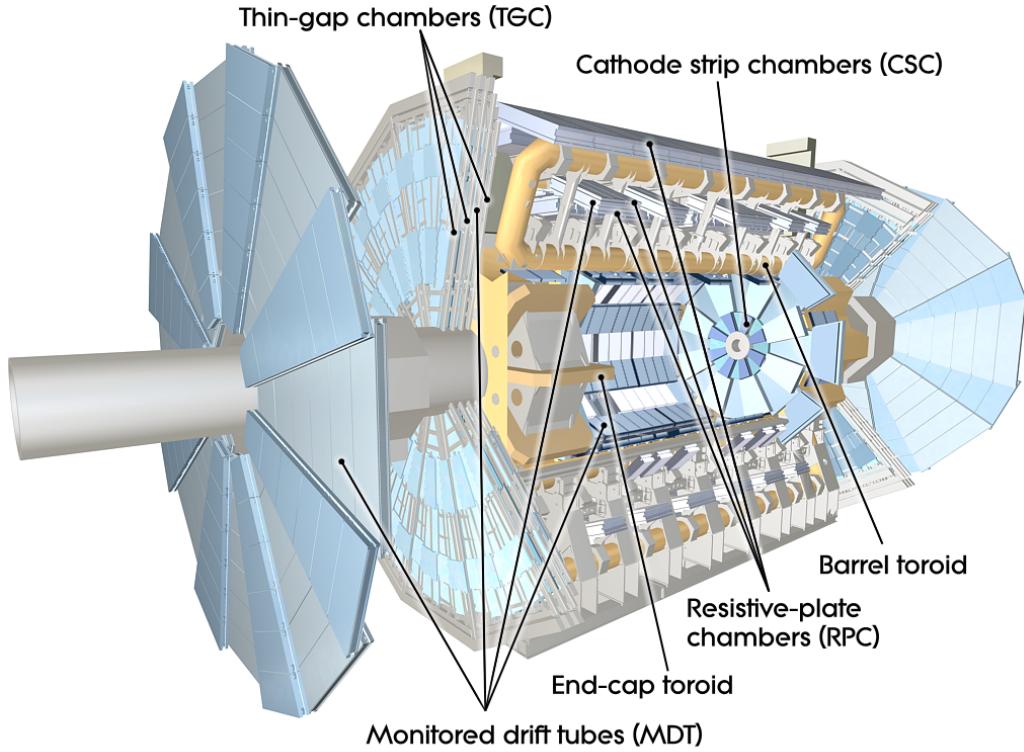
802 and the Forward LAr Calorimeter. Similar to the EM calorimeters, these are  
803 sampling calorimeters that alternate steel (dense material) with an active layer  
804 (plastic scintillator).

805 The barrel Tile calorimeter extends out to  $|\eta| < 1.7$ . There are again three layers,  
806 which combined give about 10 interactions length of distance, which provides excellent  
807 stopping power for hadrons. This is critical to avoid excess *punchthrough* to the muon  
808 spectrometer beyond the hadronic calorimeters. The first layer has a depth of 1.5  
809 interaction lengths. The second layer is again the thickest at a depth of 4.1 interaction  
810 lengths; most of the energy of incoming particle is deposited here. Both the first and  
811 second layer have segmentation of about  $\Delta\eta = 0.1, \Delta\phi = 0.1$ . Generally, one does not  
812 need as fine of granularity in the hadronic calorimeter, since the energy depositions  
813 in the hadronic calorimeters will be summed into the composite objects we know as  
814 jets. The third layer has a thickness of 1.8 interaction lengths, with a segmentation of  
815  $\Delta\eta = 0.2, \Delta\phi = 0.1$ . The use of multiple layers allows one to understand the induced  
816 hadronic shower as it propagates through the detector material.

817 The endcap LAr hadronic calorimeter covers the region  $1.5 < |\eta| < 3.2$ . It is  
818 again a sampling calorimeter; the active material is LAr with a copper absorbed. It  
819 does not use the accordion shape of the other calorimeters; it has a “standard” flat  
820 shape perpendicular to the interaction point. The segmentation varies with  $\eta$ . For  
821  $1.5 < |\eta| < 2.5$ , the cells are  $\Delta\eta = 0.1, \Delta\phi = 0.1$ ; in the region  $2.5 < |\eta| < 3.2$ , the  
822 cells are  $\Delta\eta = 0.2, \Delta\phi = 0.2$  in size.

823 The final calorimeter in ATLAS is the forward LAr calorimeter. Of those  
824 subdetectors which are used for standard reconstruction techniques, the FCal sits  
825 at the most extreme values of  $3.1 < |\eta| < 4.9$ . The FCal itself is made of three  
826 subdetectors; FCal1 is actually an electromagnetic module, while FCal2 and FCal3  
827 are hadronic. The absorber in FCal1 is copper, with a liquid argon active medium.  
828 FCal2 and FCal3 also use a liquid argon active medium, with a tungsten absorber.

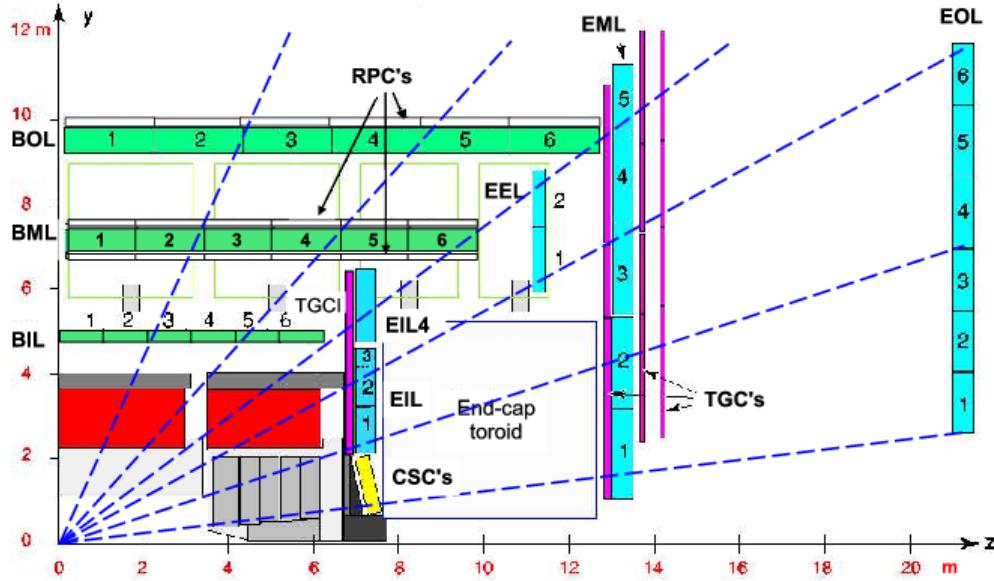
Figure 5.10: The ATLAS muon spectrometer



## 829 5.4 Muon Spectrometer

830 The muon spectrometer is the final major subdetector of the ATLAS detector.  
831 The muon spectrometer sits outside the hadronic calorimetry, with pseudorapidity  
832 coverage out to  $|\eta| < 2.7$ . The MS is a huge detector, with some detector elements  
833 existing as far as 11 m in radius from the interaction point. This system is used  
834 almost exclusively to measure the momenta of muons; these are the only measured  
835 SM particles which consistently exit the hadronic calorimeters. These systems provide  
836 a rough measurement, which is used in triggering (described in Ch.5.5), and a precise  
837 measurement to be used in offline event reconstruction as described in Ch.???. The  
838 MS produces tracks in a similar way to the ID; the hits in each subdetector are  
839 recorded and then tracks are produced from these hits. Muon spectrometer tracks are  
840 largely independent of the ID tracks due to the independent solenoidal and toroidal  
841 magnet systems used in the ID and MS respectively. The MS consists of four separate

Figure 5.11: A schematic in  $z/\eta$  showing the location of the subdetectors of the muon spectrometer

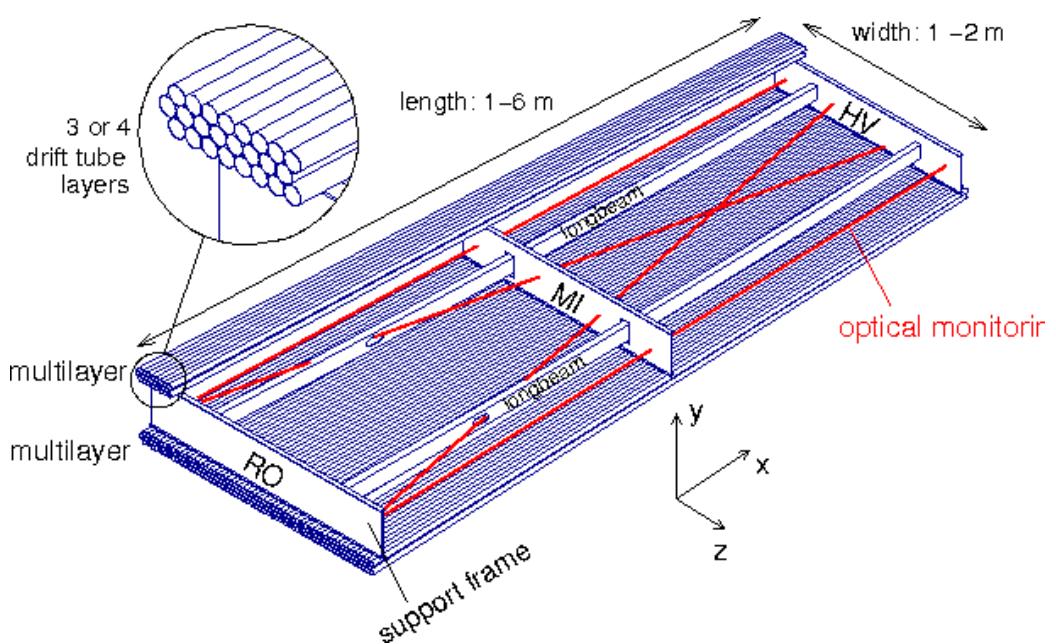


842 subdetectors: the barrel region is covered by the Resistive Plate Chambers (RPCs)  
 843 and Monitored Drift Tubes (MDTs) while the endcaps are covered by MDTs, Thin  
 844 Gap Chambers (TGCs), and Cathode Strip Chambers (CSCs).

## 845 Monitored Drift Tubes

846 The MDT system is the largest individual subdetector of the MS. MDTs provide  
 847 precision measurements of muon momenta as well as fast measurements used for  
 848 triggers. There are 1088 MDT chambers providing coverage out to pseudorapidity  
 849  $|\eta| < 2.7$ ; each consists of an aluminum tube containing an argon- $\text{CO}_2$  gas mixture.  
 850 In the center of each tube there  $50\mu\text{m}$  diameter tungsten-rhenium wire at a voltage of  
 851 3080 V. A muon entering the tube will induce ionization in the gas, which will “drift”  
 852 towards the wire due to the voltage. One measures this ionization as a current in the  
 853 wire; this current comes with a time measurement related to how long it takes the  
 854 ionization to drift to the wire.

Figure 5.12: Schematic of a Muon Drift Tube chamber



855 These tubes are layered in a pattern shown in Fig.5.12. Combining the measure-  
856 ments from the tubes in each layer gives good position resolution. The system consists  
857 of three subsystems of these layers, at 5 m, 7m, and 9 m from the interaction point.  
858 The innermost layer is directly outside the hadronic calorimeter. The combination of  
859 these three measurements gives precise momenta measurements for muons.

## 860 Resistive Plate Chambers

861 The RPC system is alternated with the MDT system in the barrel; the first two layers  
862 of RPC detectors surround the second MDT layer while the third is outside the final  
863 MDT layer. The RPC system covers pseudorapidity  $|\eta| < 1.05$ . Each RPC consists  
864 of two parallel plates at a distance of 2 mm surrounding a  $\text{C}_2\text{H}_2\text{F}_4$  mixture. The  
865 electric field between these plates is 4.9k kV/mm. Just as in the MDTs, an incoming  
866 muon ionizes the gas, and the deposited ionization is collected by the detector (in this  
867 case on the plates). It is quite fast, but with a relatively poor spatial resolution of  
868 1 cm. Still, it can provide reasonable  $\phi$  resolution due to its large distance from the  
869 interaction point. This is most useful in triggering, where the timing requirements are  
870 quite severe. The RPCs are also complement the MDTs by providing a measurement  
871 of the non-bending coordinate.

## 872 Cathode Strip Chambers

873 The CSCs are used in place of MDTs in the first layer of the endcaps. This region, at  
874  $2.0 < |\eta| < 2.7$ , has higher particle multiplicity at the close distance to the interaction  
875 point from low-energy photons and neutrons. The MDTs were not equip to deal with  
876 the higher particle rate of this region, so the CSCs were designed to deal with this  
877 deficiency.

878 Each CSC consists multiwire proportional chambers, oriented radially outward  
879 from the interaction point. These chambers overlap partially in  $\phi$ . The wires contain

Figure 5.13: Photo of the installation of Cathode Strip Chambers and Monitored Drift Tubes



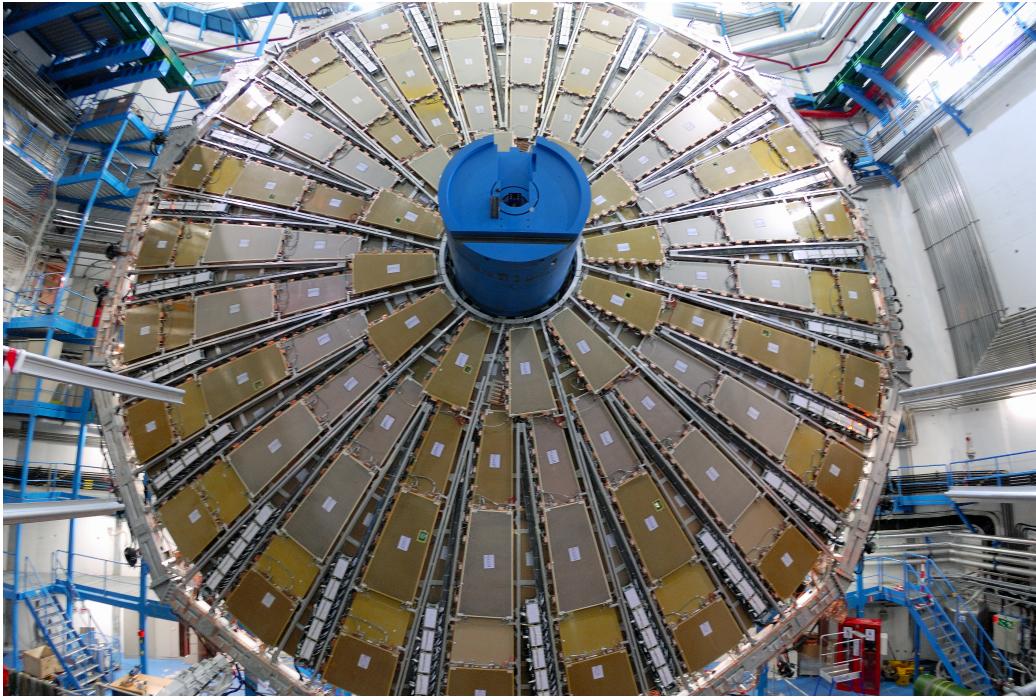
880 a gas mixture of argon and CO<sub>2</sub>, which is ionized when muons enter. The detectors  
881 operate with a voltage of 1900 V, with much lower drift times than the MDTs. They  
882 provide less hits than MDTs, but their lower drift times lower uptime and reduce the  
883 amount of detector overload.

884 The CSCs are arranged into four planes on the wheels of the muon spectrometer,  
885 as seen in Fig.???. There are 32 CSCs in total, with 16 on each side of the detector  
886 in  $\eta$ .

### 887 **Thin Gap Chambers**

888 The TGCs serve the purpose of the RPCs in the endcap at pseudorapidity of  $1.05 <$   
889  $|\eta| < 2.4$ ; they provide fast measurements used in triggering. The TGCs are also  
890 multiwire proportional chambers a la the CSCs. The fast readouts necessary for  
891 trigger are provided by a high electric field and a small wire-to-wire distance of 1.8  
892 mm. These detectors provide both  $\eta$  and  $\phi$  information, allowing the trigger to use  
893 as much information as possible when selecting events.

Figure 5.14: Photo of a muon Big Wheel, consisting of Thin Gap Chambers



## 894 5.5 Trigger System

895 The data rate delivered by the LHC is staggering [89]. In the 2016 dataset, the  
896 collision rate was 40 MHz, meaning a *bunch spacing* of 25 ns. In each of the event,  
897 as we saw in Ch.??, there are many proton-proton collisions. Most of the collisions  
898 are uninteresting, such as elastic scattering of protons, or even inelastic scattering  
899 leading to low-energy dijet events. These types of events have been studied in detail  
900 in previous experiments.

901 Even if one is genuinely interested in these events, it's *impossible* to save all of  
902 the information available in each event. If all events were written "to tape" (as the  
903 jargon goes), ATLAS would store terabytes of data per second. We are limited to only  
904 about 1000 Hz readout by computing processing time and storage space. We thus  
905 implement a *trigger* which provides fast inspection of events to drastically reduce  
906 the data rate from the 40 MHz provided by the LHC to the 1000 Hz we can write to  
907 tape for further analysis.

908       The ATLAS trigger system consists of a two-level trigger, known as the Level-  
909       1 trigger (L1 trigger) and the High-Level Trigger (HLT)<sup>4</sup>. Trigger selections are  
910       organized into *trigger chains*, where events passing a particular L1 trigger are passed  
911       to a corresponding HLT trigger. For example, one would require a particular high- $p_T$   
912       muon at L1, with additional quality requirements at HLT. One can also use HLT  
913       triggers as prerequisites for each other, as is done in some triggers requiring both jets  
914       and  $E_T^{\text{miss}}$ .

915       **Level-1 Trigger**

916       The L1 trigger is hardware-based, and provides the very fast rejection needed to  
917       quickly select events of interest. The L1 trigger uses only what is known as *prompt*  
918       data to quickly identify interesting events. Only the calorimeters and the triggering  
919       detectors (RPCs and TGCs) of the MS are fast enough to be considered at L1,  
920       since the tracking reconstruction algorithms used by the ID and the more precise  
921       MS detectors are very slow. This allows quick identification of events with the  
922       most interesting physical objects : large missing transverse momentum and high-  
923        $p_T$  electrons, muons, and jets.

924       L1 trigger processing is done locally. This means that events are selected without  
925       considering the entire available event. Energy deposits over some threshold are  
926       reconstructed as *regions of interest*. These RoIs are then compared using pattern  
927       recognition hardware to “expected” patterns for the given RoIs. Events with RoIs  
928       matching these expected patterns are then handed to the HLT through the Central  
929       Trigger Processor. This step alone lowers the data rate down by about three orders  
930       of magnitude.

---

<sup>4</sup>In Run1, ATLAS ran with a three-level trigger system. The L1 was essentially as today; the HLT consisted of two separate systems known as the L2 trigger and the Event Filter (EF). This was changed to the simpler system used today during the shutdown between Run1 and Run2.

931 **High-Level Trigger**

932 The HLT performs the next step, taking the incoming data rate from the L1 trigger  
933 of  $\sim 75$  kHz down to the  $\sim 1$  kHz that can be written to tape. The HLT really  
934 performs much like a simplified offline reconstruction, using many common quality  
935 and analysis cuts to eliminate uninteresting events. This is done by using computing  
936 farms located close to the detector, which process events in parallel. Individually, each  
937 event which enters the computing farms takes about 4 seconds to reconstruct; the  
938 HLT reconstruction time also has a long tail, which necessitates careful monitoring  
939 of the HLT to ensure smooth operation.

940 HLT triggers are targetted to a particular physics process, such as a  $E_T^{\text{miss}}$  trigger,  
941 single muon trigger, or multijet trigger. The collection of all triggers is known as  
942 the trigger *menu*. Since many low-energy particles are produced in collisions, it is  
943 necessary to set a *trigger threshold* on the object of interest; this is really just a fancy  
944 naming for a trigger  $p_T$  cut. Due to the changing luminosity conditions of the LHC,  
945 these thresholds change constantly, mostly by increasing thresholds with increasing  
946 instantaneous luminosity. This allows an approximately constant number of events to be  
947 written for further analysis. Triggers which have rates higher than those designated  
948 by the menu are *prescaled*. This means writing only some fraction of the triggered  
949 events. Of course, for physics analyses, one wishes to investigate all data events  
950 passing some set of analysis cuts, so often one uses the “lowest threshold unprescaled  
951 trigger”. *Turn-on curves* allow one to select the needed offline analysis cut to ensure  
952 the trigger is fully efficient. An example turn-on curve for the  $E_T^{\text{miss}}$  triggers used in  
953 the signal region of this analysis is shown in ??.

954 The full set of the lowest threshold unprescaled triggers considered here can be  
955 found in Table 5.1. These are the lowest unprescaled triggers associated to the SUSY  
956 signal models and Standard Model backgrounds considered in this thesis. More  
957 information can be found in [89].

Table 5.1: High-Level Triggers used in this thesis. Descriptions of loose, medium, tight, and isolated can be found in [89]. The  $d_0$  cut refers to a quality cut on the vertex position; this was removed from many triggers in 2016 to increase sensitivity to displaced vertex signals. For most triggers, the increased thresholds in 2016 compared to 2015 were designed to keep the rate approximately equal. The exception is the  $E_T^{\text{miss}}$  triggers; see 5.5.

Physics Object	Trigger	$p_T$ (GeV)	Threshold	Level-1 Seed	Additional Requirements	Approximate Rate (Hz)
<b>2015 Data</b>						
$E_T^{\text{miss}}$	HLT_xe70	70		L1_XE50	-	60
	HLT_mu24_iloose_L1 <b>M145</b>			L1_MU15	isolated, loose	130
Muon	HLT_mu50	50		L1_MU15	-	30
Muon	HLT_e24_1hmedium_l1 <b>B4se_L1EM20VH</b>			L1_EM20VH	medium OR isolated, loose	140
Electron	HLT_e60_1hmedium	60		L1_EM20VH	medium	10
Electron	HLT_e120_1hloose	120		L1_EM20VH	loose	<10
Electron	HLT_g120_loose	120		L1_EM20VH	loose	20
<b>2016 Data</b>						
$E_T^{\text{miss}}$	HLT_xe100_mht_L1 <b>XE500</b>			L1_XE50	-	180
	HLT_mu24_ivarmedium <b>4</b>			L1_MU20	medium	120
Muon	HLT_mu50	50		L1_MU20	-	40
Muon	HLT_e24_l1tight_no <b>d4ivarloose</b>			L1_EM22VHT	tight with no $d_0$ or loose	110
Electron	HLT_e60_1hmedium_no <b>d0</b>			L1_EM22VHT	medium with no $d_0$	10
Electron	HLT_e140_1hloose_no <b>d0</b>			L1_EM22VHT	loose with no $d_0$	<10
Electron	HLT_g140_loose	140		L1_EM22VHT	loose	20

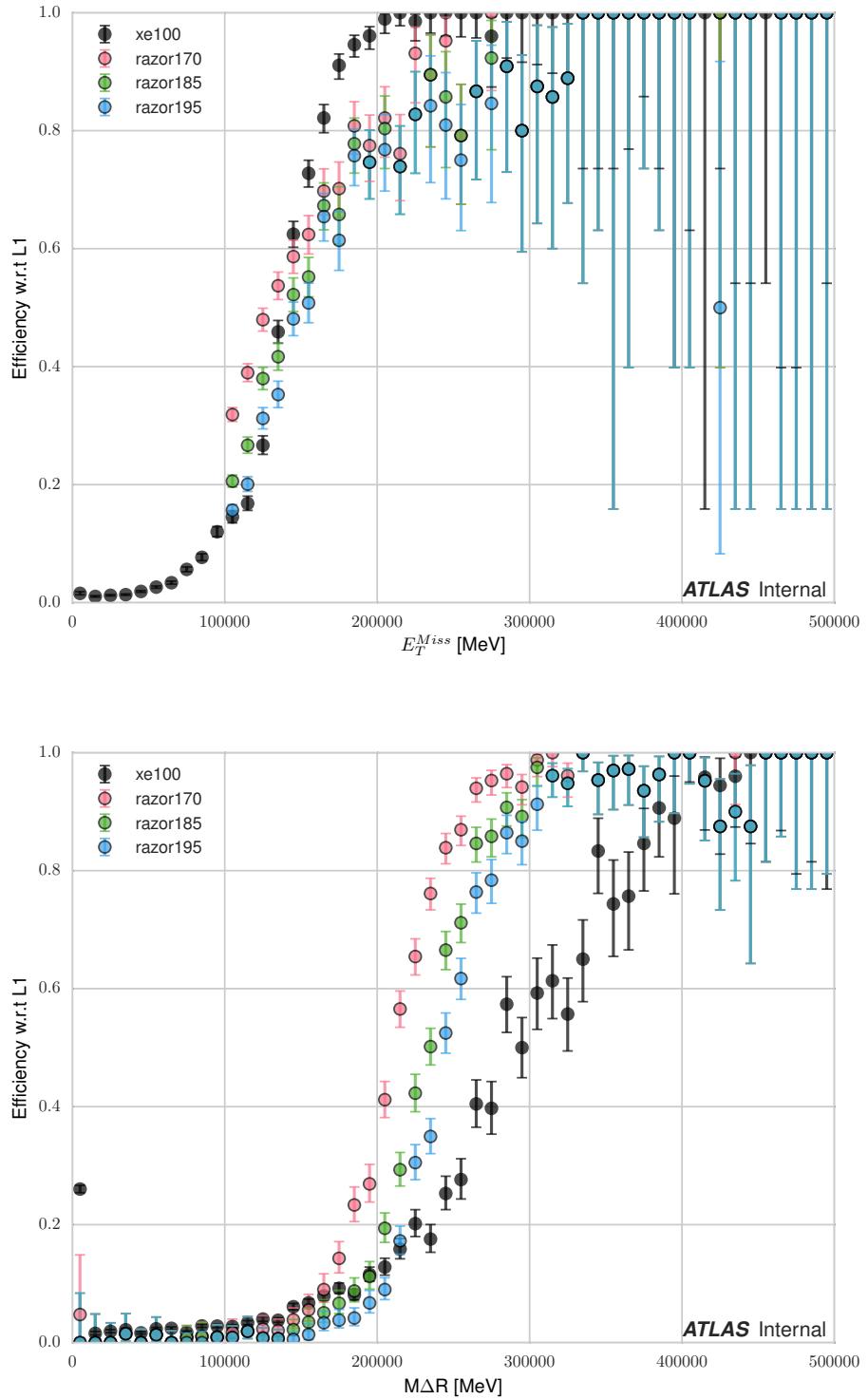
958 **Razor Triggers**

959 For the analysis presented in this thesis, the *razor triggers* were developed. These are  
960 topological triggers, combining both jet and  $E_T^{\text{miss}}$  information to select interesting  
961 events. In particular, they use the razor variable  $M_{\Delta}^R$  which will be described in  
962 Chapter ??.

963 Based on 2015 run conditions, these triggers would have allowed the use of a lower  
964 offline  $E_T^{\text{miss}}$  cut with a similar rate to the nominal  $E_T^{\text{miss}}$  triggers. This can be seen  
965 in the turn-on curves shown in Figure 5.15. The razor triggers are fully efficient at  
966 nearly 100 GeV lower than the corresponding  $E_T^{\text{miss}}$  triggers in  $M_{\Delta}^R$ .

967 There was a quite big change in the 2016 menu, which increased the rate given to  
968  $E_T^{\text{miss}}$  triggers drastically. This can be seen in the difference in rate shown between  
969  $E_T^{\text{miss}}$  triggers in 2015 and 2016 in Table 5.1. This allowed the  $E_T^{\text{miss}}$  triggers to  
970 maintain a lower threshold throughout the dataset used in this thesis.

Figure 5.15: Turn-on curves for the razor triggers and nominal  $E_T^{\text{miss}}$  trigger. The razor triggers show a much sharper turn-on in  $M_{\Delta}^R$  relative to the  $E_T^{\text{miss}}$  trigger. The converse is true for the  $E_T^{\text{miss}}$  triggers.





971

## Chapter 6

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972

### Event Reconstruction

973 This chapter describes the reconstruction algorithms used within ATLAS. We will  
974 make the distinction between the “primitive” objects which are reconstructed from  
975 the detector signals from the “composite” physics objects we use in measurements  
976 and searches for new physics.

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#### 977 6.1 Primitive Object Reconstruction

978 The primitive objects reconstructed by ATLAS are *tracks* and (calorimeter) *clusters*.  
979 These are reconstructed directly from tracking hits and calorimeter energy deposits  
980 into cells. Tracks can be further divided into inner detector and muon spectrom-  
981 eter tracks. Calorimeter clusters can be divided into sliding-window clusters and  
982 topological clusters (topoclusters).

#### 983 Inner Detector Tracks

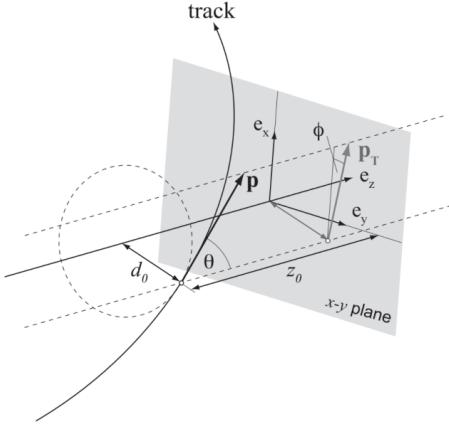
984

985 Inner detector tracks are reconstructed from hits in the inner detector. These hits  
986 indicate that a charged particle has passed through the detector material. Due to the  
987 2 T solenoid in the inner detector, the hits associated with any individual particle will  
988 be curved; this allows one to measure the momentum of the particle. In any given  
989 event, there is upwards of , making it impossible to do any sort of combinatorics to  
990 reconstruct tracks<sup>1</sup>. There are two algorithms used by ATLAS track reconstruction,

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Figure 6.1: The parameters associated to a track.



991 known as *inside-out* and *outside-in*.

992 ATLAS first employs the inside-out algorithm. First, one assumes the track begins  
 993 at the interaction point. Moving out from the interaction point, one creates track  
 994 seeds. Track seeds are proto-tracks constructed from three hits; these hits can be  
 995 distributed as three pixel hits, two pixel hits and one SCT hit, or three SCT hits.

site

996 One extrapolates the track and uses a [combinatorial Kalman filter](#), which adds the  
 997 rest of the pixel and SCT hits to the seeds. This is done seed by seed, so it avoids  
 998 the combinatorial complexity involved with checking all hits with all seeds. At this  
 999 point, the algorithm applies an additional filter to avoid ambiguities from nearby  
 1000 tracks. The TRT hits are then added to the seeds in the same procedure; in this way,  
 1001 all hits are associated to a track.

1002 The next step is to figure out the correct kinematics of the track. This is  
 1003 done by applying a fitting algorithm which outputs the best-fit track parameters  
 1004 by minimizing the track distance from hits, weighted by each hit's resolution. These  
 1005 parameters are  $(d_0, z_0, \eta, \phi, q/p)$  where  $d_0$  ( $z_0$ ) is the transverse (longitudinal) impact  
 1006 parameter and  $q/p$  is the charge over the track momenta. This set of parameters  
 1007 uniquely defines the trajectory of the charged particle associated to the track; an  
 1008 illustration of a track with these parameters is shown in Fig.6.1.

1009       The other track reconstruction algorithm is the outside-in algorithm. As the name  
1010      implies, in this case, we start from the outside of the inner detector, in the TRT, and  
1011      extend the tracks in. One begins by seeding from TRT hits, and extending the track  
1012      back towards the center of the detector. The same fitting procedure is used as in  
1013      the inside-out algorithm to find the optimal track parameters. This algorithm is  
1014      particularly important for finding tracks which originate from interactions with the  
1015      detector material, especially the SCT. For tracks from primary vertices, this often  
1016      finds the same tracks as the inside-out algorithm, providing an important check on  
1017      the consistency of the tracking procedure.

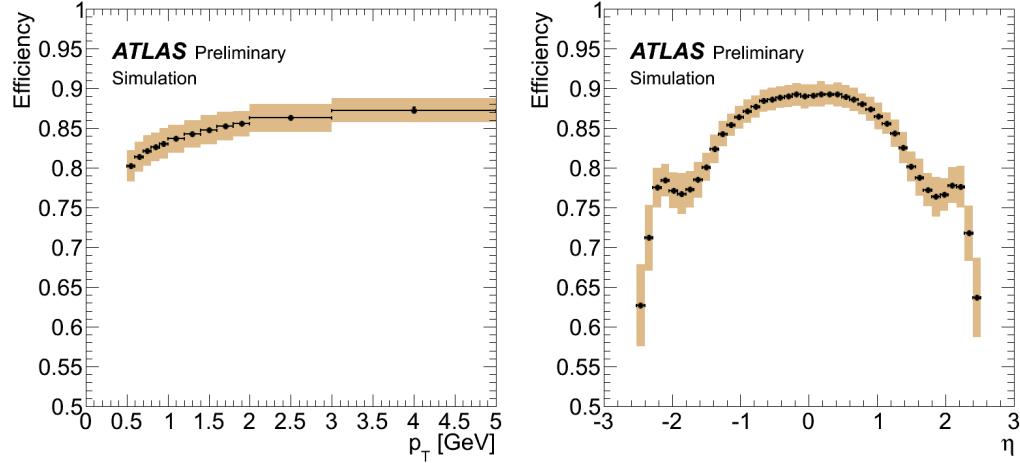
1018       In the high luminosity environment of the LHC, even the tracks reconstructed  
1019      from precision detectors such as those of ATLAS inner detector can sometimes lead  
1020      to fake tracks from simple combinatoric chance. Several quality checks are imposed  
1021      after track fitting which reduce this background. Seven silicon (pixel + SCT) hits  
1022      are required for all tracks. No more than two holes are allowed in the pixel detector;  
1023      holes are expected measurements from the track that are missing in the pixel detector.  
1024      Finally, tracks with poor fit quality, as measured by  $\chi^2/ndf$ , are also rejected. Due  
1025      to the high quality of the silicon measurements in the pixel detector and SCT, these  
1026      requirements give good track reconstruction efficiency, as seen in Fig.6.2 for simulated  
1027      events[[ATL-COM-PHYS-2012-1541](#)].

## 1028 Sliding-window clusters

1029       The sliding-window algorithm is a way to combine calorimeter cells into composite  
1030      objects (clusters) to be used as inputs for other algorithms[[90](#)]. Sliding-window  
1031      clusters are the primary inputs to electron and photon reconstruction, as described  
1032      below. As described in Ch.??, the electromagnetic calorimeter has high granularity,  
1033      with a cell size of  $(\eta, \phi) = (.025, .025)$  in the coarsest second layer throughout most  
1034      of the calorimeter. The “window” consists of 3 by 5 cells in the  $(\eta, \phi)$  space; all

Figure 6.2: Track reconstruction efficiency as a function of track  $p_T$  and  $\eta$ . The efficiency is defined as the number of reconstructed tracks divided by the number of generate charged particles.

(a) Track reconstruction as a function of  $p_T$ . (b) Track reconstruction as a function of  $\eta$ .



1035 layers are added on this same 2D space. One translates this window over the space  
 1036 and seeds a cluster whenever the energy sum of the cells is maximized. If the seed  
 1037 energy is greater than 2.5 GeV, this seed is called a sliding-window cluster. This  
 1038 choice was motivated to optimize the reconstruction efficiency of proto-electrons and  
 1039 proto-photons while rejecting fakes from electronic noise and additional particles from  
 1040 pileup vertices.

## 1041 Topological clusters

1042 Topoclusters are the output of the algorithm used within ATLAS to combine  
 1043 hadronic and electromagnetic calorimeter cells in a way which extracts signal from  
 1044 a background of significant electronic noise[91]. They are the primary input to the  
 1045 algorithms which reconstruct jets.

1046 Topological clusters are reconstructed from calorimeter cells in the following way.  
 1047 First, one maps all cells onto a single  $\eta - \phi$  plane so one can speak of *neighboring*  
 1048 cells. Two cells are considered neighboring if they are in the same layer and directly

1049 adjacent, or if they are in adjacent layers and overlap in  $\eta - \phi$  space. The *significance*  
1050  $\xi_{\text{cell}}$  of a cell during a given event is

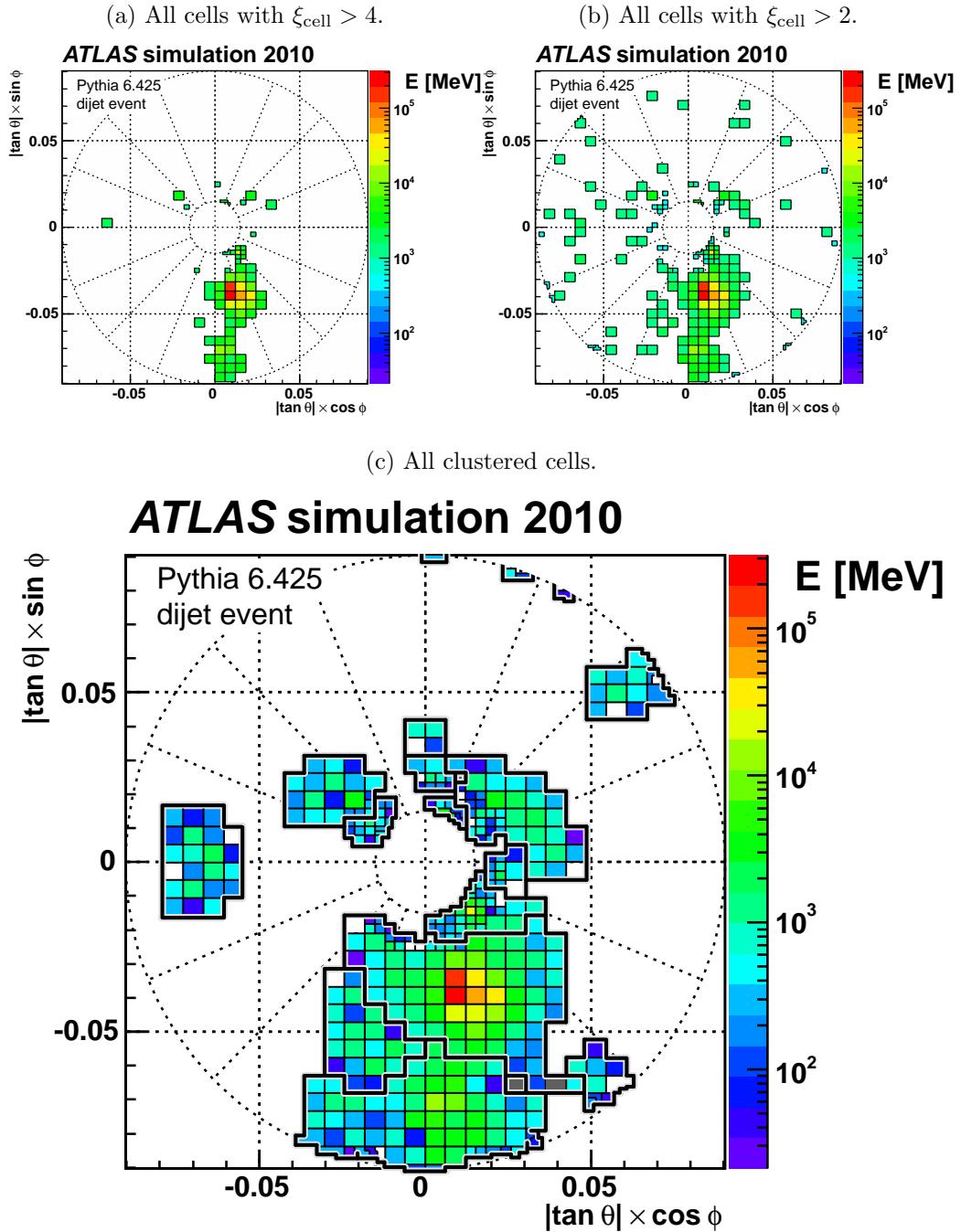
$$\xi_{\text{cell}} = \frac{E_{\text{cell}}}{\sigma_{\text{noise,cell}}} \quad (6.1)$$

1051 where  $\sigma_{\text{noise,cell}}$  is measured for each cell in ATLAS and  $E_{\text{cell}}$  measures the current  
1052 energy level of the cell. One thinks of this as the measurement of the energy *over*  
1053 *threshold* for the cell.

1054 Topocluster *seeds* are defined as calorimeter cells which have a significance  $\xi_{\text{cell}} >$   
1055 4. These are the inputs to the algorithm; one iteratively tests all cells adjacent to these  
1056 seeds for  $\xi_{\text{cell}} > 2$ . Each cells passing this selection is then added to the topocluster,  
1057 and the procedure is repeated. When the algorithm reaches the point where there  
1058 are no additional adjacent cells with  $\xi_{\text{cell}} > 2$ , every positive-energy cell adjacent to  
1059 the current proto-cluster is added. This collection of cells is summed; the summed  
1060 object is known as a topocluster. An example of this procedure for a simulation dijet  
1061 event is shown in Fig.6.3.

1062 There are two calibrations used for clusters. These are known as the electromag-  
1063 cite  
1064 netic (EM) scale and the local cluster weighting (LCW) scale. The EM scale is the  
1065 energy read directly out of the calorimeters as described. This scale is appropriate for  
1066 electromagnetic processes. The LCW scale applies additional scaling to the clusters  
1067 based on the shower development. This allows the cluster energy to be corrected for  
1068 calorimeter non-compensation and the differences in the hadronic and electromagnetic  
1069 calorimeters' responses. This scale provides additional corrections that improve the  
1070 accuracy of hadronic energy measurements. This thesis only uses the EM scale  
1071 corrections unfortunately; LCW scaling requires additional measurements that only  
1072 became available with additional data. Due to the jet calibration procedure that  
1073 we will describe below, it is also a relatively complicated procedure to rederive the  
“correct” jet energy.

Figure 6.3: Example of topoclustering on a simulated dijet event.



1074 **Muon Spectrometer Tracks**

1075 Muon spectrometer tracks are fit using the same algorithms as the ID tracks, but  
1076 different subdetectors. The tracks are seeded by hits in the MDTs or CSCs. After  
1077 seeding in the MDTs and CSCs, the hits from all subsystems are refit as the final  
1078 MS track. These tracks are used as inputs to the muon reconstruction, as we will see  
1079 below.

1080 **6.2 Physics Object Reconstruction and Quality**  
1081 **Identification**

1082 There are essentially six objects used in ATLAS searches for new physics: electrons,  
1083 photons, muons,  $\tau$ -jets, jets, and  $E_T^{\text{miss}}$ . The reconstruction of these objects is  
1084 described here; in this thesis,  $\tau$  lepton jets are not treated differently from other  
1085 hadronic jets. A very convenient summary plot is shown in Fig.6.4. This process  
1086 produces candidate objects, which are then identified by quality.

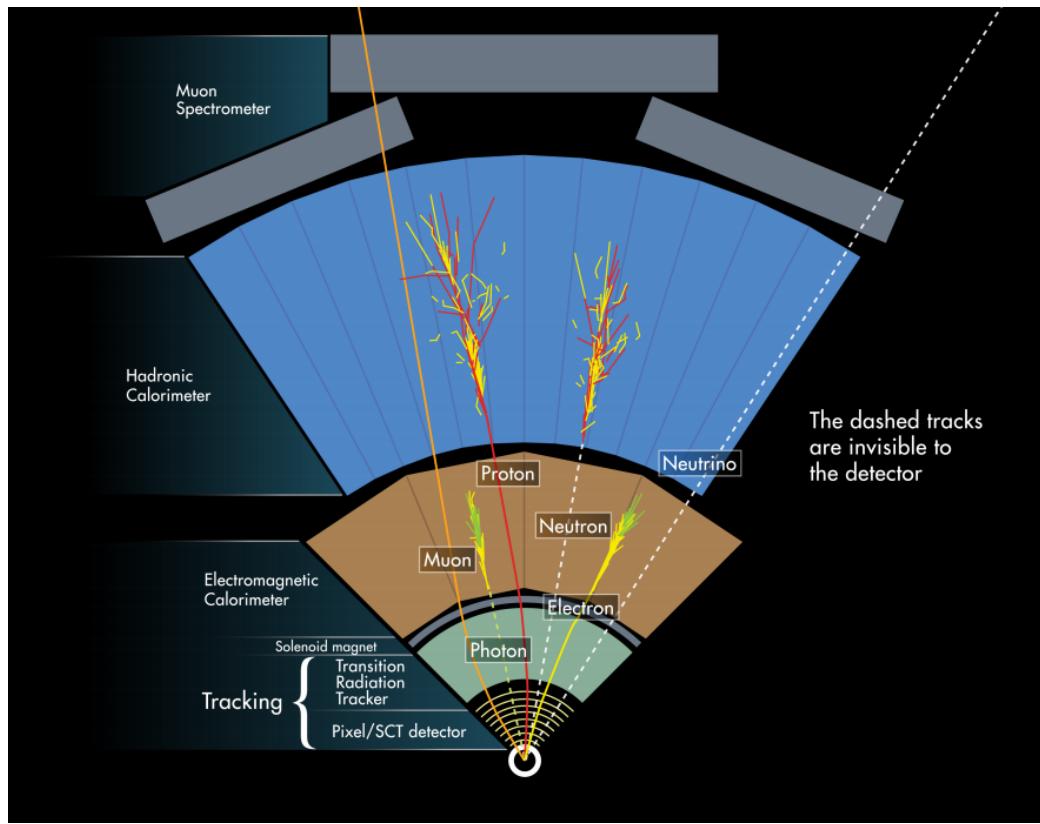
1087 One often wishes to understand “how certain” we are that a particular object  
1088 is truly the underlying physics object. In ATLAS, we often generically consider, in  
1089 order, *very loose*, *loose*, *medium*, and *tight* objects<sup>2</sup>. These are ordered in terms of  
1090 decreasing object efficiency, or equivalently, decreasing numbers of fake objects. We  
1091 will also describe briefly the classification of objects into these categories.

1092 In this thesis, we present a search for new physics in a zero lepton final state; we  
1093 will provide additional details about jet and  $E_T^{\text{miss}}$  reconstruction.

---

<sup>2</sup> These are not all used for all objects, but it’s conceptually useful to think of these different categories.

Figure 6.4: The interactions of particles with the ATLAS detector. Solid lines indicate the particle is interacting with the detector, while dashed lines are shown where the particle does not interact.



## 1094 Electrons and Photons

### 1095 Reconstruction

1096 The reconstruction of electrons and photons (often for brevity called “electromagnetic  
 1097 objects”) is very similar [PERF-2013-03 , 92, 93]. This is because the reconstruction  
 1098 begins with the energy deposit in the calorimeter in the form of an electromagnetic  
 1099 shower. For any incoming  $e/\gamma$ , this induces many more electrons and photons in the  
 1100 shower; the measurement in the calorimeter is similar for these two objects.

1101 One thus begins the reconstruction of electromagnetic objects from the sliding-  
 1102 window clusters reconstructed from the EM calorimeter, as described in Sec.6.1.  
 1103 These  $E > 2.5$  GeV clusters provide the primary seed for electrons and photons. One

then looks for all ID tracks within  $\Delta R < 0.3$ . We “match” the track and cluster if

check delta<sup>104</sup>  
R defined  
somewhere

1105 they are within  $\Delta\phi < 0.2$  in the direction of track curvature, or  $\Delta\phi < 0.05$  in the  
1106 direction opposite the track curvature. Those track-cluster seeds with tracks pointing  
1107 to the primary vertex are reconstructed as electrons.

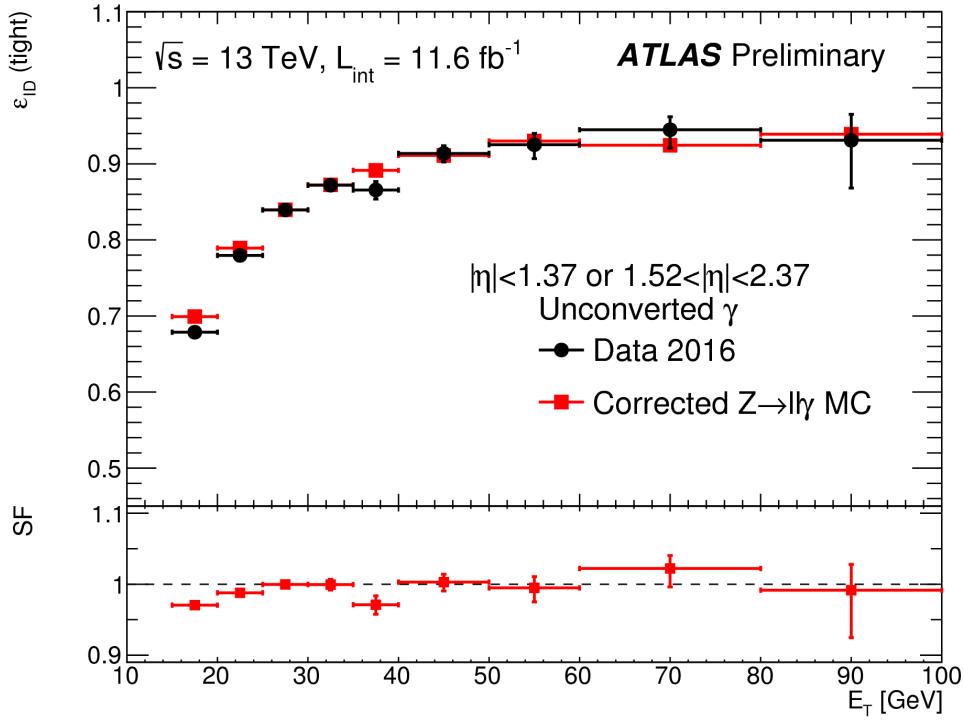
1108 For photons, we have two options to consider, known as *converted* and *unconverted*  
1109 photons. Due to the high energy of the LHC collisions, typical photons have energy  
1110  $\sim 1$  GeV; at this scale, photons interact almost exclusively via pair-production in the  
1111 presence of the detector material . If the track-cluster seed has a track which does  DIAGRAM  
1112 not point at the primary vertex, we reconstruct this object as a converted photon.  
1113 This happens since the photon travels a distance before decay into two electrons, and  
1114 see the tracks coming from this secondary vertex. Those clusters which do not have  
1115 any associated tracks are then reconstructed as an unconverted photon.

1116 The final step in electromagnetic object reconstruction is the final energy value  
1117 assigned to these objects; this process is different between electrons and photons due  
1118 to their differing signatures in the EM calorimeter. In the barrel, electrons energies  
1119 are assigned as the sum of the 3 clusters in  $\eta$  and 7 clusters in  $\phi$  to account for the  
1120 electron curving in the  $\phi$  direction. Barrel photons are assigned the energy sum of  
1121 (3, 5) clusters in  $(\eta, \phi)$  space. In the endcap, the effect of the magnetic field on the  
1122 electrons is smaller, and there is a coarser granularity. Both objects sum the (5, 5)  
1123 clusters for their final energy value.

## 1124 Quality Identification

1125 Electrons have a number of important backgrounds which can give fakes. Fake  
1126 electrons come primarily from secondary vertices in hadron decays or misidentified  
1127 hadronic jets. To reduce these backgrounds, quality requirements are imposed on  
1128 electron candidates. Loose electrons have requirements imposed on the shower  
1129 shapes in the electromagnetic calorimeter and on the quality of the associated ID  
1130 track. There is also a requirement that there is a small energy deposition in the

Figure 6.5: Unconverted photon efficiency as measured in [ATL-PHYS-PUB-2016-015].



hadronic calorimeter behind the electron, to avoid jets being misidentified as electrons (low hadronic leakage). Medium and tight electrons have increasingly stronger requirements on these variables, and additional requirements on the isolation (as measured by  $\Delta R$ ) and matching of the ID track momentum and the calorimeter energy deposit.

Photons are relatively straightforward to measure, since there are few background processes[ATL-PHYS-PUB-2016-015]. The primary one is pion decays to two photons, which can cause a jet to be misidentified as photon. Loose photons have requirements on the shower shape and hadronic leakage. Tight photons have tighter shower shape cuts, especially on the high granularity first layer of the EM calorimeter. The efficiency for unconverted tight photons as a function of  $p_T$  is shown in

1142 **Muons**

1143 **Reconstruction**

1144 Muons are reconstructed using measurements from all levels of the ATLAS detec-  
1145 tor[94]. They leave a ID track, a small, characteristic deposition in the EM calorime-  
1146 ter, and then a track in the muon spectrometer. The primary reconstruction technique  
1147 produces a so-called *combined* muon. “Combined” means using a combination of the  
1148 ID and MS tracks to produce the final reconstructed muon kinematics. This is done  
1149 by refitting the hits associated to both tracks, and using this refit track for the muon  
1150 kinematics. This process produces the best measured muons, although several other  
1151 worse algorithms are used when the full detector information is missing. An example  
1152 is in the region  $2.5 < |\eta| < 2.7$  outside the ID acceptance; in this region, MS tracks  
1153 are used without the corresponding ID tracks.

1154 **Quality Identification**

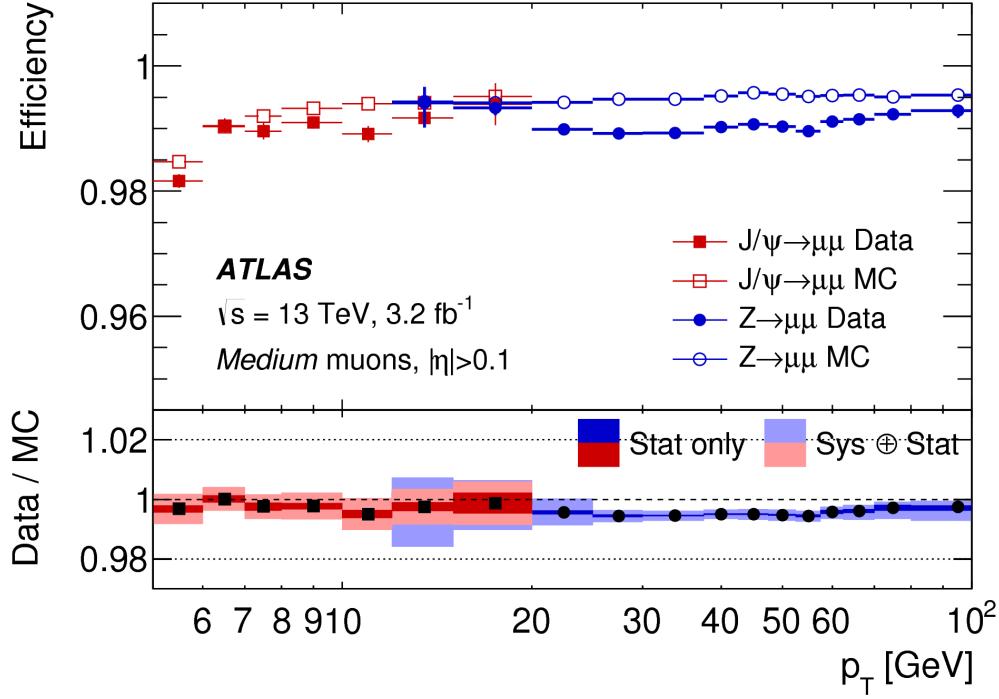
Several additional criteria are used to assure muon measurements are free of significant background contributions, especially from pion and kaon decays to muons. Muons produced via these decay processes are often characterized by a “kink”. Candidate muons with a poor fit quality, characterized by  $\chi^2/\text{n.d.f.}$ , are thus rejected. Additionally, the absolute difference in momentum measurements between the ID and MS provide another handle, since the other decay products from hadron decays carry away some amount of the initial hadron momentum. This is measured by

$$\rho' = \frac{|p_T^{\text{ID}} - p_T^{\text{MS}}|}{p_T^{\text{Combined}}}. \quad (6.2)$$

Additionally, there is a requirement on the  $q/p$  significance, defined as

$$S_{q/p} = \frac{|(q/p)^{\text{ID}} - (q/p)^{\text{MS}}|}{\sqrt{\sigma_{\text{ID}}^2 + \sigma_{\text{MS}}^2}}. \quad (6.3)$$

Figure 6.6: Medium muon efficiency as measured in [94].



1155 The  $\sigma_{ID,MS}$  in the denominator of Eq.6.3 are the uncertainties on the corresponding  
 1156 quantity from the numerator. Finally, cuts are placed on the number of hits in the  
 1157 various detector elements.

1158 Subsequently tighter cuts on these variables allow one to define the different muon  
 1159 identification criteria. Loose muons have the highest reconstruction efficiency, but  
 1160 the highest number of fake muons, since there are no requirements on the number  
 1161 of subdetector hits and the loosest requirements on the suite of quality variables.  
 1162 Medium muons consist of Loose muons with tighter cuts on the quality variables;  
 1163 they also require more than three MDT hits in at least two MDT layers. These are  
 1164 the default used by ATLAS analyses. Tight muons have stronger cuts than those of  
 1165 the medium selection, and reducing the reconstruction efficiency. The reconstruction  
 1166 efficiency as a function of  $p_T$  can be seen for Medium muons in Fig.6.6.

1167 **Jets**

1168

cite pa-  
per/note

1169 Jets are composite objects corresponding to many physical particles. This is a  
1170 striking difference from the earlier particles. Fortunately, we normally (and in this  
1171 thesis) care about the original particle produced in primary collision. In the SM, this  
1172 corresponds to quarks and gluons; due to the hadronization process described in 2, free  
1173 quarks and gluons spontaneously hadronize and produce a hadronic shower, which  
1174 we call a jet. These showers can be measured by the EM and hadronic calorimeters,  
1175 and the charged portions can be measured in the ID. The first question is how to  
1176 combine these measurements into a composite object representing the underlying  
1177 physical parton. This is done via jet algorithms.

1178 **Jet Algorithms**

1179 It might seem straightforward to combine the underlying physical particles into a  
1180 jet. There are three important characteristics required for any jet reconstruction  
1181 algorithm to be used by ATLAS.

- 1182     • Collinear safety - if any particle with four-vector  $p$  is replaced by two particles  
1183         of  $p_1, p_2$  with  $p = p_1 + p_2$ , the subsequent jet should not change
- 1184     • Radiative (infrared) safety - if any particle with four-vector  $p$  radiates a particle  
1185         of energy  $\alpha \rightarrow 0$ , the subsequent jet should not change
- 1186     • Fast - the jet algorithm should be “fast enough” to be useable by ATLAS  
1187         computing resources

1188 The first two requirements can be seen in terms of requirements on soft gluon emission.  
1189 Since partons emit arbitrarily soft gluons freely, one should expect the algorithms

1190 to not be affected by this emission. The final requirement is of course a practical  
1191 limitation.

The algorithms in use by ATLAS (and CMS) which satisfies these requirements are collectively known as the  $k_T$ algorithms [95–97]. These algorithms iteratively combine the “closest” objects, defined using the following distance measures :

$$d_{ij} = \min(k_{T,i}^{2p}, k_{T,j}^{2p}) \frac{\Delta_{ij}^2}{R^2} \quad (6.4)$$
$$d_{iB} = k_{Ti}^{2p}$$

1192 In Eq.6.4,  $k_T$ ,  $i$  is the transverse momentum of  $i$ -th jet *constituent*,  $\Delta_{ij}$  is the angular  
1193 distance between the constituents. Both  $R$  and  $p$  are adjustable parameters;  $R$  is  
1194 known as the (jet) *cone size* and  $p$  regulates the power of the energy vs the geometrical  
1195 scales. The algorithm sequence, for a given set of objects  $i$  with four-vector  $k$  :

- 1196 1. Find the minimum distance in the set of all  $d_{ij}$  and  $d_{iB}$ .
- 1197 2. If the distance is one of the  $d_{ij}$ , combine the input pair of object  $i, j$  and return  
1198 to (1). If the distance is one of the  $d_{iB}$ , remove the object from the list, call it  
1199 a jet, and return to (1).

1200 This process ends when all objects  $i$  have been added to a jet.

1201 Any choice of  $(p, R)$  has the requirements of collinear and radiative safety. In  
1202 essence, the choice is then to optimize based on speed and the potential for new  
1203 physics discoveries. In ATLAS, we make the choice of  $p = -1$ ; this is also known as  
1204 the *anti- $k_T$* algorithm. The choice of  $R = 0.4$  is used for the distance parameter of  
1205 the jets.

1206 The primary “nice” quality of this algorithm can be seen with the following  
1207 example. Consider three inputs to an anti- $k_T$ algorithm, all with  $\eta = 0$  :

- 1208 • Object 1 :  $(p_T, \phi) = (30 \text{ GeV}, 0)$
- 1209 • Object 2 :  $(p_T, \phi) = (20 \text{ GeV}, -0.2)$

1210 • Object 3 :  $(p_T, \phi) = (10 \text{ GeV}, 0.2)$

1211 • Object 4 :  $(p_T, \phi) = (1 \text{ GeV}, 0.5)$

1212 . In the case shown, it seems natural to first combine the “bigger” objects 1 and 2.  
1213 These then pick up the extra small object 3, and object 4 is not included in the jet.  
1214 This is exactly what is done by the anti- $k_T$  algorithm. The (normal)  $k_T$  algorithm with  
1215  $p = 1$  instead combines the smallest objects, 3 and 4, first. Object 1 and 2 combine  
1216 to form their own jet, instead of these jets picking up object 3. This behavior is not  
1217 ideal due to the effects of pileup, as we will see in the next section.

## 1218 Jet Reconstruction

1219 In ATLAS, jets are reconstructed using multiple different objects as inputs, including  
1220 tracks, “truth” objects, calorimeter clusters, and *particle flow objects* (PFOs) . For  
1221 physics analyses, ATLAS primarily uses jets reconstructed from calorimeter clusters,  
1222 but we will describe the others here, as they are often used for derivations of  
1223 systematic uncertainties or future prospects.

cite pflow  
paper

1224 Calorimeter jets are reconstructed using topoclusters using the anti- $k_T$  algorithm  
1225 with  $R = 0.4$ . The jet reconstruction algorithm is run on the collection of all  
1226 topoclusters reconstructed as in Sec.6.1. Both EM and LCW scale clusters are used  
1227 in the ATLAS reconstruction software and produce two sets of jets for analysis. As  
1228 stated above, this thesis presents an analysis using jets reconstructed using EM scale  
1229 clusters; we refer to these as *EM jets*.

1230 Tracks can be used as inputs to jet reconstruction algorithms. Jets reconstructed  
1231 from tracks are known as *track jets*. Since the ID tracks do not measure neutral  
1232 objects, these jets measure an incorrect energy. However, these are still useful for  
1233 checks and derivations of systematic uncertainties.

1234     *Truth* jets are reconstructed from *truth* particles. In this case, truth is jargon  
1235 for simulation; in simulation, the actual simulated particles are available and used as  
1236 inputs to the jet reconstruction algorithms. Similarly to track jets, these are not useful  
1237 in and of themselves. Instead, truth jets are used for comparisons and derivations of  
1238 systematic uncertainties.

1239     The last object generally used as inputs to jet reconstruction algorithms are  
1240 particle flow objects (PFOs). Particle flow objects are reconstructed by associating  
1241 tracks and clusters through a combination of angular distance measures and detector  
1242 response measurements to create a composite object which contains information from  
1243 both the ID and the calorimeters. For calorimeter clusters which do not have any  
1244 associated ID track, the cluster is simply the PFO. The natural association between  
1245 tracks and clusters provides easy pileup subtraction since tracks are easily associated  
1246 to the primary vertex. This technique is generally used in CMS, and ATLAS has  
1247 been slow to adopt the same. As pileup has increased, the utility of using PFOs as  
1248 inputs to jet reconstruction has increased as well.

## 1249 Jet Calibration

1250     Jets as described in the last section are still *uncalibrated*. Even correcting the cluster  
1251 energies using the LCW does not fully correct the jet energy, due to particles losing  
1252 energy in the calorimeters. The solution to this is the *jet energy scale* (JES). The  
1253 JES is a series of calibrations which on average restore the correct truth jet energy  
1254 for a given reconstructed jet. These steps are shown in Fig.6.7 and described here.

1255     The first step is the origin correction; this adjusts the jet to point at the primary  
1256 vertex. Next, is the jet-area based pileup correction. This step subtracts the  
1257 “average” pileup as measured by the energy density  $\rho$  outside of the jets and assumes  
1258 this is a good approximation for the pileup inside the jet. One then removes energy  
1259  $\Delta E = \rho \times A_{\text{jet}}$  in this step. The residual pileup correction makes a final offset

1260 correction by parametrizing the change in jet energy as a function of the number of  
1261 primary vertices  $N_{\text{PV}}$  and the average number of interactions  $\mu$ .

1262 The next step is the most important single correction; it is known as the  
1263 AbsoluteEtaJES step. Due to the use of non-compensation and sampling calorimeters  
1264 in ATLAS, the measured energy of a jet is a fraction of the true energy of the outgoing  
1265 parton. Additionally, due to the use of different technologies and calorimeters  
1266 throughout the detector, there are directional biases induced by these effects. The  
1267 correction bins a multiplicative factor in  $p_{\text{T}}$  and  $\eta$  which scales the reconstructed jets  
1268 to corresponding truth jet  $p_{\text{T}}$ . This step does not entirely correct the jets, since it is  
1269 entirely a simulation-based approach.

1270 The final steps are known as the global sequential calibration (GSC) and the  
1271 residual in-situ calibration. The GSC uses information about the jet showering shape  
1272 to apply additional corrections based on the expected shape of gluon or quark jets.  
1273 The final step is the residual in-situ calibration, which is only applied to data. This  
1274 step uses well-measured objects recoiling off a jet to provide a final correction to the  
1275 jets in data. In the low  $p_{\text{T}}$  region ( $20 \text{ GeV} \sim < p_{T,\text{jet}} \sim < 200 \text{ GeV}$ ),  $Z \rightarrow ll$  events are  
1276 used as a reference object. In the middle  $p_{\text{T}}$  region ( $100 \text{ GeV} \sim < p_{T,\text{jet}} \sim < 600 \text{ GeV}$ ),  
1277 the reference object is a photon, while in the high  $p_{\text{T}}$  region ( $p_{T,\text{jet}} \sim > 200 \text{ GeV}$ ), the  
1278 high  $p_{\text{T}}$  jet is compared to multiple smaller  $p_{\text{T}}$  jets; the reference object is this group  
1279 of multijets. After this final correction, the data and MC scales are identical up to  
1280 the corresponding uncertainties; the combined JES uncertainty as a function of  $p_{\text{T}}$  is  
1281 shown in Fig.??.

## 1282 **B-jets**

## 1283 **Missing Transverse Momentum**

1284

Figure 6.7: The steps used by ATLAS to calibrate jets

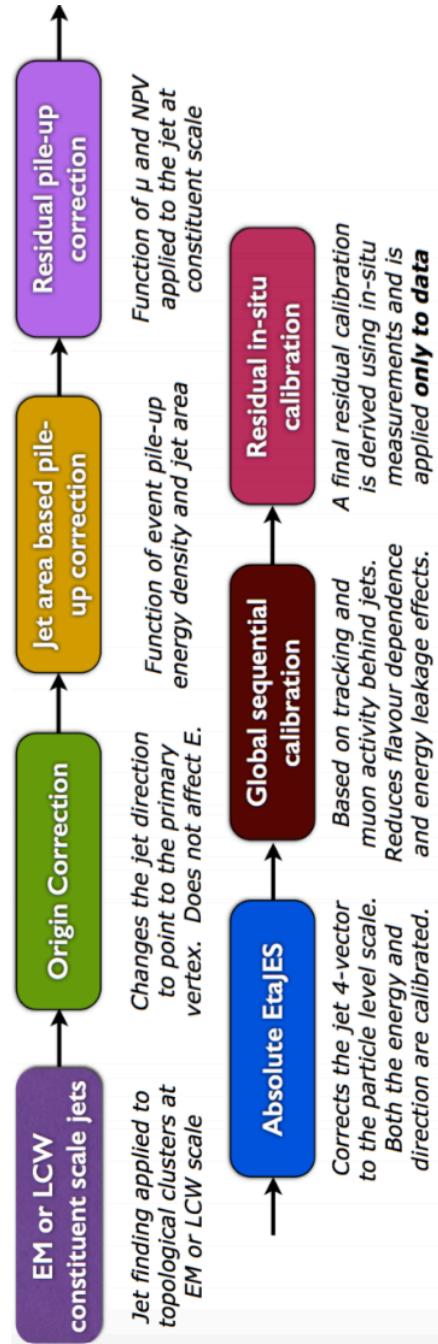
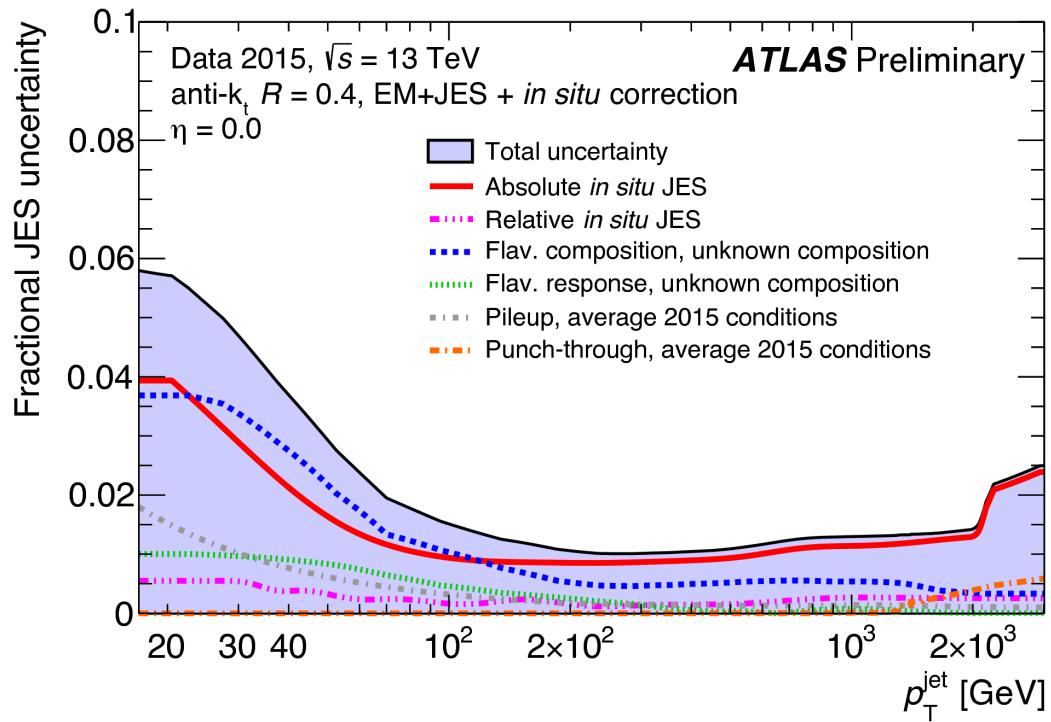


Figure 6.8: Combined jet energy scale uncertainty as a function of  $p_T$  at  $\eta = 0$ .



1285 **6.3 Maybe PFlow?**



1286

## Chapter 7

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1287

### *The Recursive Jigsaw Technique*

1288 Here you can write some introductory remarks about your chapter. I like to give each  
1289 sentence its own line.

1290 When you need a new paragraph, just skip an extra line.

1291 **7.1 Razor variables**

1292 By using the asterisk to start a new section, I keep the section from appearing in the  
1293 table of contents. If you want your sections to be numbered and to appear in the  
1294 table of contents, remove the asterisk.

1295 **7.2 SuperRazor variables**

1296 **7.3 The Recursive Jigsaw Technique**

1297 **7.4 Variables used in the search for zero lepton**

1298 **SUSY**



1299

## Chapter 8

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1300

### *Title of Chapter 1*



1301

## Chapter 9

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1302

### *Title of Chapter 1*

1303 Here you can write some introductory remarks about your chapter. I like to give each  
1304 sentence its own line.

1305 When you need a new paragraph, just skip an extra line.

1306 **9.1 Object reconstruction**

1307 **Photons, Muons, and Electrons**

1308 **Jets**

1309 **Missing transverse momentum**

1310 Probably longer, show some plots from the PUB note that we worked on

<sub>1311</sub> **9.2 Signal regions**

<sub>1312</sub> **Gluino signal regions**

<sub>1313</sub> **Squark signal regions**

<sub>1314</sub> **Compressed signal regions**

<sub>1315</sub> **9.3 Background estimation**

<sub>1316</sub> **Z vv**

<sub>1317</sub> **W ev**

<sub>1318</sub> **ttbar**

1319

## Chapter 10

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1320

### *Title of Chapter 1*

1321 Here you can write some introductory remarks about your chapter. I like to give each  
1322 sentence its own line.

1323 When you need a new paragraph, just skip an extra line.

## 1324 **10.1 Statistical Analysis**

1325 maybe to be moved to an appendix

## 1326 **10.2 Signal Region distributions**

## 1327 **10.3 Pull Plots**

## 1328 **10.4 Systematic Uncertainties**

## 1329 **10.5 Exclusion plots**



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1330

## *Conclusion*

1331 Here you can write some introductory remarks about your chapter. I like to give each  
1332 sentence its own line.

1333 When you need a new paragraph, just skip an extra line.

1334 **10.6 New Section**

1335 By using the asterisk to start a new section, I keep the section from appearing in the  
1336 table of contents. If you want your sections to be numbered and to appear in the  
1337 table of contents, remove the asterisk.



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1598 In this appendix, we provide a brief overview of the basic ingredients involved in  
 1599 construction of the Standard Model Lagrangian : quantum field theory, symmetries,  
 1600 and symmetry breaking.

1601 **Quantum Field Theory**

1602

1603 In this section, we provide a brief overview of the necessary concepts from  
 1604 Quantum Field Theory (QFT).

cite Yuval's  
lectures  
and notes  
somehow

1605 In modern physics, the laws of nature are described by the “action”  $S$ , with the  
 1606 imposition of the principle of minimum action. The action is the integral over the cite  
 1607 spacetime coordinates of the “Lagrangian density”  $\mathcal{L}$ , or Lagrangian for short. The  
 1608 Lagrangian is a function of “fields”; general fields will be called  $\phi(x^\mu)$ , where the  
 1609 indices  $\mu$  run over the space-time coordinates. We can then write the action  $S$  as

$$S = \int d^4x \mathcal{L}[\phi_i(x^\mu), \partial_\mu \phi_i(x^\mu)] \quad (10.1)$$

1610 where we have an additional summation over  $i$  (of the different fields). Generally,  
 1611 we impose the following constraints on the Lagrangian :

- 1612 1. Translational invariance - The Lagrangian is only a function of the fields  $\phi$  and  
 1613 their derivatives  $\partial_\mu \phi$
- 1614 2. Locality - The Lagrangian is only a function of one point  $x_\mu$  in spacetime.

1615     3. Reality condition - The Lagrangian is real to conserve probability.

1616     4. Lorentz invariance - The Lagrangian is invariant under the Poincarégroup of  
1617       spacetime.

1618     5. Analyticity - The Lagrangian is an analytical function of the fields; this is to  
1619       allow the use of perturbation theory.

1620     6. Invariance and Naturalness - The Lagrangian is invariant under some internal  
1621       symmetry groups; in fact, the Lagrangian will have *all* terms allowed by the  
1622       imposed symmetry groups.  
maybe add 622  
in ref here

1623     7. Renormalizability - The Lagrangian will be renormalizable - in practice, this  
1624       means there will not be terms with more than power 4 in the fields.

1625     The key item from the point of view of this thesis is that of “Invariance and  
1626       Natural”. We impose a set of “symmetries” and then our Lagrangian is the most  
1627       general which is allowed by those symmetries.

## 1628     **Symmetries**

1629     Symmetries can be seen as the fundamental guiding concept of modern physics.

cite?     1630     Symmetries are described by “groups”. To illustrate the importance of symmetries  
1631       and their mathematical description, groups, we start here with two of the simplest  
1632       and most useful examples :  $\mathbb{Z}_2$  and  $U(1)$ .

### 1633      $\mathbb{Z}_2$ symmetry

1634      $\mathbb{Z}_2$  symmetry is the simplest example of a “discrete” symmetry. Consider the most  
1635       general Lagrangian of a single real scalar field  $\phi(x_\mu)$

$$\mathcal{L}_\phi = \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - \frac{m^2}{2}\phi^2 - \frac{\mu}{2\sqrt{2}}\phi^3 - \lambda\phi^4 \quad (10.2)$$

Now we *impose* the symmetry

$$\mathcal{L}(\phi) = \mathcal{L}(-\phi) \quad (10.3)$$

1636 This has the effect of restricting the allowed terms of the Lagrangian. In particular,  
 1637 we can see the term  $\phi^3 \rightarrow -\phi^3$  under the symmetry transformation, and thus must  
 1638 be disallowed by this symmetry. This means under the imposition of this particular  
 1639 symmetry, our Lagrangian should be rewritten as

$$\mathcal{L}_\phi = \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - \frac{m^2}{2}\phi^2 - \lambda\phi^4 \quad (10.4)$$

1640 The effect of this symmetry is that the total number of  $\phi$  particles can only change  
 1641 by even numbers, since the only interaction term  $\lambda\phi^4$  is an even power of the field.  
 1642 This symmetry is often imposed in supersymmetric theories, as we will see in Chapter  
 1643 3.

## 1644 **$U(1)$ symmetry**

1645  $U(1)$  is the simplest example of a continuous (or *Lie*) group. Now consider a theory  
 1646 with a single complex scalar field  $\phi = \text{Re } \phi + i \text{Im } \phi$

$$\mathcal{L}_\phi = \delta_{i,j} \frac{1}{2}\partial_\mu\phi_i\partial^\mu\phi_j - \frac{m^2}{2}\phi_i\phi_j - \frac{\mu}{2\sqrt{2}}\phi_i\phi_j\phi_k\phi_l - \lambda\phi_i\phi_j\phi_k\phi_l \quad (10.5)$$

1647 where  $i, j, k, l = \text{Re}, \text{Im}$ . In this case, we impose the following  $U(1)$  symmetry  
 1648 :  $\phi \rightarrow e^{i\theta}, \phi^* \rightarrow e^{-i\theta}$ . We see immediately that this again disallows the third-order  
 1649 terms, and we can write a theory of a complex scalar field with  $U(1)$  symmetry as

$$\mathcal{L}_\phi = \partial_\mu\phi\partial^\mu\phi^* - \frac{m^2}{2}\phi\phi^* - \lambda(\phi\phi^*)^2 \quad (10.6)$$

## 1650 Local symmetries

1651 The two examples considered above are “global” symmetries in the sense that the  
1652 symmetry transformation does not depend on the spacetime coordinate  $x_\mu$ . We know  
1653 to look at local symmetries; in this case, for example with a local  $U(1)$  symmetry, the  
1654 transformation has the form  $\phi(x_\mu) \rightarrow e^{i\theta(x_\mu)}\phi(x_\mu)$ . These symmetries are also known  
1655 as “gauge” symmetries; all symmetries of the Standard Model are gauge symmetries.

There are wide-ranging consequences to the imposition of local symmetries. To begin, we note that the derivative terms of the Lagrangian 10.2 are *not* invariant under a local symmetry transformation

$$\partial_\mu \phi(x_\mu) \rightarrow \partial_\mu(e^{i\theta(x_\mu)}\phi(x_\mu)) = (1 + i\theta(x_\mu))e^{i\theta(x_\mu)}\phi(x_\mu) \quad (10.7)$$

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1657 This leads us to note that the kinetic terms of the Lagrangian are also not invariant  
1658 under a gauge symmetry. This would lead to a model with no dynamics, which is  
1659 clearly unsatisfactory.

1660 Let us take inspiration from the case of global symmetries. We need to define a  
1661 so-called “covariant” derivative  $D^\mu$  such that

$$D^\mu \phi \rightarrow e^{iq\theta(x^\mu)D^\mu}\phi \quad (10.8)$$

$$D^\mu \phi^* \rightarrow e^{-iq\theta(x^\mu)D^\mu}\phi^* \quad (10.9)$$

$$(10.10)$$

1662 Since  $\phi$  and  $\phi^*$  transforms with the opposite phase, this will lead the invariance  
1663 of the Lagrangian under our local gauge transformation. This  $D^\mu$  is of the following  
1664 form

$$D^\mu = \partial_\mu - igqA^\mu \quad (10.11)$$

1665 where  $A^\mu$  is a vector field we introduce with the transformation law

$$A^\mu \rightarrow A^\mu - \frac{1}{g} \partial_\mu \theta \quad (10.12)$$

1666 and  $g$  is the coupling constant associated to vector field. This vector field  $A^\mu$  is  
1667 also known as a “gauge” field.

1668 Since we need to add all allowed terms to the Lagrangian, we define

$$F^{\mu\nu} = A^\mu A^\nu - A^\nu A^\mu \quad (10.13)$$

1669 and then we must also add the kinetic term :

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4} F^{\mu\nu} F_{\mu\nu} \quad (10.14)$$

1670 The most general renormalizable Lagrangian with fermion and scalar fields can  
1671 be written in the following form

$$\mathcal{L} = \mathcal{L}_{\text{kin}} + \mathcal{L}_\phi + \mathcal{L}_\psi + \mathcal{L}_{\text{Yukawa}} \quad (10.15)$$

## 1672 Symmetry breaking and the Higgs mechanism

1673 Here we view some examples of symmetry breaking. We investigate breaking of a  
1674 global  $U(1)$  symmetry and a local  $U(1)$  symmetry. The SM will break the electroweak  
1675 symmetry  $SU(2)xU(1)$ , and in Chapter 3 we will see how supersymmetry must also  
1676 be broken.

1677 There are two ideas of symmetry breaking

- 1678 • Explicit symmetry breaking by a small parameter - in this case, we have a small  
1679 parameter which breaks an “approximate” symmetry of our Lagrangian. An  
1680 example would be the theory of the single scalar field 10.2, when  $\mu \ll m^2$  and

$\mu \ll \lambda$ . In this case, we can often ignore the small term when considering low-energy processes.

- Spontaneous symmetry breaking (SSB) - spontaneous symmetry breaking occurs when the Lagrangian is symmetric with respect to a given symmetry transformation, but the ground state of the theory is *not* symmetric with respect to that transformation. This can have some fascinating consequences, as we will see in the following examples

Symmetry breaking a

### U(1) global symmetry breaking

Consider the theory of a complex scalar field under the  $U(1)$  symmetry, or the transformation

$$\phi \rightarrow e^{i\theta} \phi \quad (10.16)$$

The Lagrangian for this theory is

$$\mathcal{L} = \partial^\mu \phi^\dagger \partial_\mu \phi + \frac{\mu^2}{2} \phi^\dagger \phi + \frac{\lambda}{4} (\phi^\dagger \phi)^2 \quad (10.17)$$

Let us write this theory in terms of two scalar fields,  $h$  and  $\xi$  :  $\phi = (h + i\xi)/\sqrt{2}$ .

The Lagrangian can then be written as

$$\mathcal{L} = \partial^\mu h \partial_\mu h + \partial^\mu \xi d\mu \xi - \frac{\mu^2}{2} (h^2 + \xi^2) - \frac{\lambda}{4} (h^2 + \xi^2)^2 \quad (10.18)$$

First, note that the theory is only stable when  $\lambda > 0$ . To understand the effect of SSB, we now enforce that  $\mu^2 < 0$ , and define  $v^2 = -\mu^2/\lambda$ . We can then write the scalar potential of this theory as :

$$V(\phi) = \lambda(\phi^\dagger \phi - v^2/2)^2 \quad (10.19)$$

Minimizing this equation with respect to  $\phi$ , we can see that the “vacuum expectation value” of the theory is

$$2 < \phi^\dagger \phi > = < h^2 + \xi^2 > = v^2 \quad (10.20)$$

1690        We now reach the “breaking” point of this procedure. In the  $(h, \xi)$  plane, the  
 1691        minima form a circle of radius  $v$ . We are free to choose any of these minima to expand  
 1692        our Lagrangian around; the physics is not affected by this choice. For convenience,  
 1693        choose  $\langle h \rangle = v, \langle \xi^2 \rangle = 0$ .

Now, let us define  $h' = h - v, \xi' = \xi$  with VEVs  $\langle h' \rangle = 0, \langle \xi' \rangle = 0$ . We can  
 then write our spontaneously broken Lagrangian in the form

$$\mathcal{L} = \frac{1}{2} \partial_\mu h' \partial^\mu h' + \frac{1}{2} \partial_\mu \xi' \partial^\mu \xi' - \lambda v^2 h'^2 - \lambda v h' (h'^2 + \xi'^2) - \lambda (h'^2 + \xi'^2)^2 \quad (10.21)$$