

1 A search for sparticles in zero lepton final states

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3 Submitted in partial fulfillment of the

4 requirements for the degree of

5 Doctor of Philosophy

6 in the Graduate School of Arts and Sciences

7 COLUMBIA UNIVERSITY

8 2016

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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*

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

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

84 Despite its impressive range of described phenomena, the Standard Model has  
 85 some theoretical and experimental deficiencies. The SM contains 26 free parameters  
 86 <sup>1</sup>. It would be more theoretically pleasing to understand these free parameters in  
 87 terms of a more fundamental theory. The major theoretical concern of the Standard  
 88 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  $\alpha_{force}$  ).

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

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

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

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

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



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**133 2.1 Overview**

134 A Standard Model is another name for a theory of the internal symmetry group  
 135  $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ , with its associated set of parameters. *The Standard*  
 136 Model refers specifically to a Standard Model with the proper parameters to describe  
 137 the universe. The SM is the culmination of years of work in both theoretical  
 138 and experimental particle physics. In this thesis, we take the view that theorists cite

139 construct a model with the field content and symmetries as inputs, and write down the  
 140 most general Lagrangian consistent with those symmetries. Assuming this model is  
 141 compatible with nature (in particular, the predictions of the model are consistent with  
 142 previous experiments), experimentalists are responsible measuring the parameters of  
 143 this model. This will be applicable for this chapter and the following one.

144 Additional theoretical background is in 9.6. The philosophy and notations are  
 145 inspired by [48, 49].

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**146 2.2 Field Content**

The Standard Model field content is

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

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

150 We observed that  $Q_L, U_R$ , and  $D_R$  are triplets under  $SU(3)_C$ ; these are the *quark*  
151 fields. The *color* group,  $SU(3)_C$  is mediated by the *gluon* field  $G^\mu(8, 1)_0$ , which has  
152 8 degrees of freedom. The fermion fields  $L_L(1, 2)_{-1}$  and  $E_R(1, 1)_{-2}$  are singlets under  
153  $SU(3)_C$ ; we call them the *lepton* fields.

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

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

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

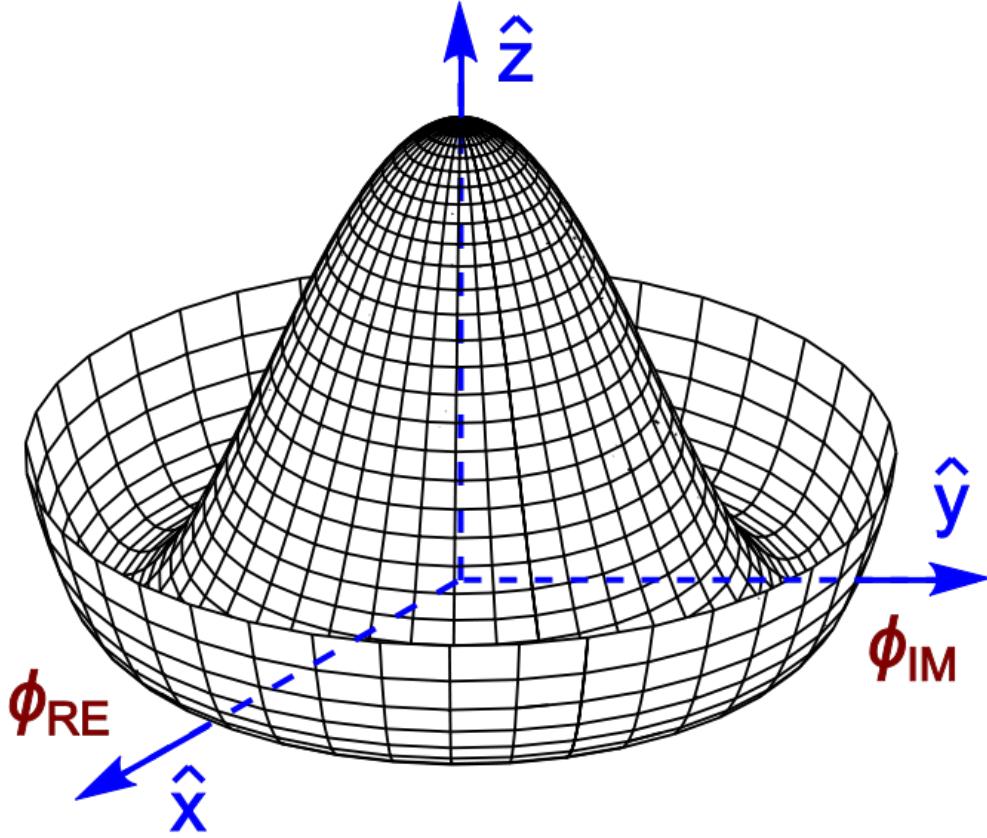
## 165 Electroweak sector

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

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

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

167      The terms in the Lagrangian 2.2 proportional to  $\mu^2$  and  $\lambda$  make up the “Higgs  
 168 potential” [50]. As normal (see Appendix 9.6), we restrict  $\lambda > 0$  to guarantee our  
 169 potential is bounded from below, and we also require  $\mu^2 < 0$ , which gives us the  
 170 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].

## 177 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

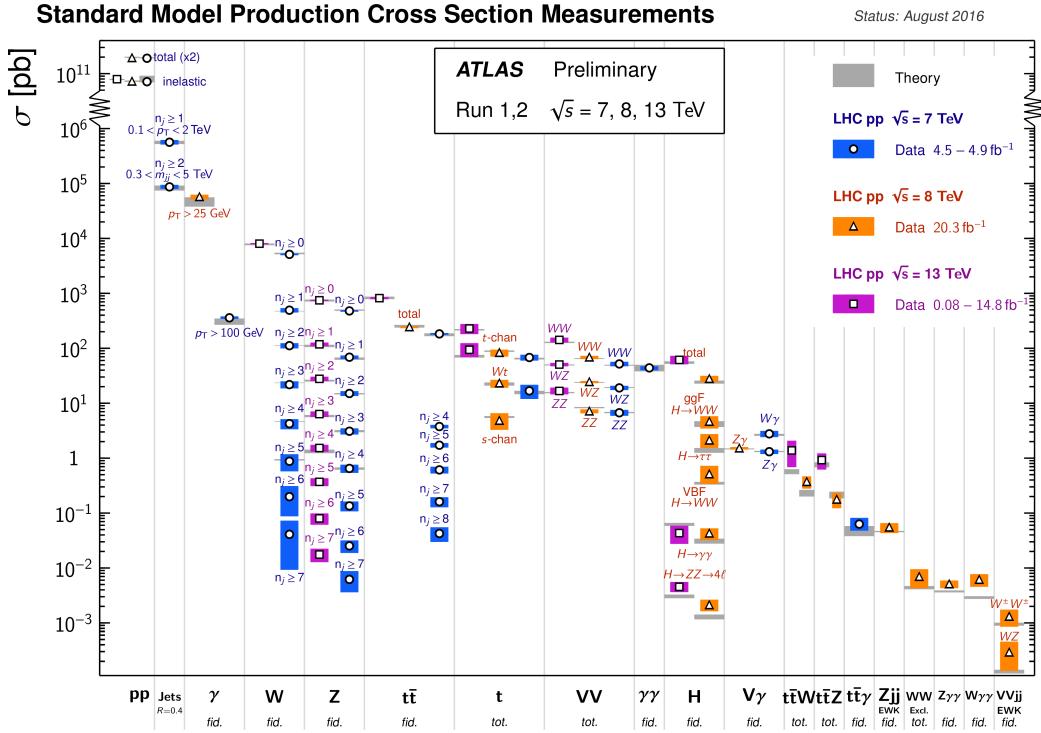
$$G_a^{\mu\nu} = \partial^\mu G_a^\nu - \partial^\nu G_a^\mu - g_s f^{abc} G_b^\mu G_c^\nu, a, b, c = 1, \dots, 8 \quad (2.13)$$

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

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

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

Figure 2.2: Cross-sections of various Standard Model processes



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

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

216 **Fermions**

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

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

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

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

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

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

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

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

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

236     Each up-like quark has charge  $q_{up} = 2/3$ , while the down-like quarks have  $q_{down} =$   
237  $-1/3$ . At the high energies of the LHC, one often makes the distinction between  
238 the light quarks ( $u, d, c, s$ ), the bottom quark, and top quark. In general, due to  
239 the hadronization process described above, the light quarks, with masses  $m_q < \sim$   
240  $1.5\text{GeV}$  are indistinguishable by LHC experiments. Their hadronic decay products  
241 generally have long lifetimes and they are reconstructed as jets.<sup>1</sup>. The bottom quark  
242 hadronizes primarily through the  $B$ -mesons, which generally travels a short distance  
243 before decaying to other hadrons. This allows one to distinguish decays via  $b$ -quarks  
244 from other jets; this procedure is known as *b-tagging* and will be discussed more in  
245 Ch.5. Due to its large mass, the top quark decays before it can hadronize; there  
246 are no bound states associated to the top quark. The top is of particular interest at  
247 the LHC; it has a striking signature through its most common decay mode  $t \rightarrow Wb$ .  
248 Decays via tops, especially  $t\bar{t}$  are frequently an important signal decay mode, or an  
249 important background process.

## 250 **Interactions in the Standard Model**

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

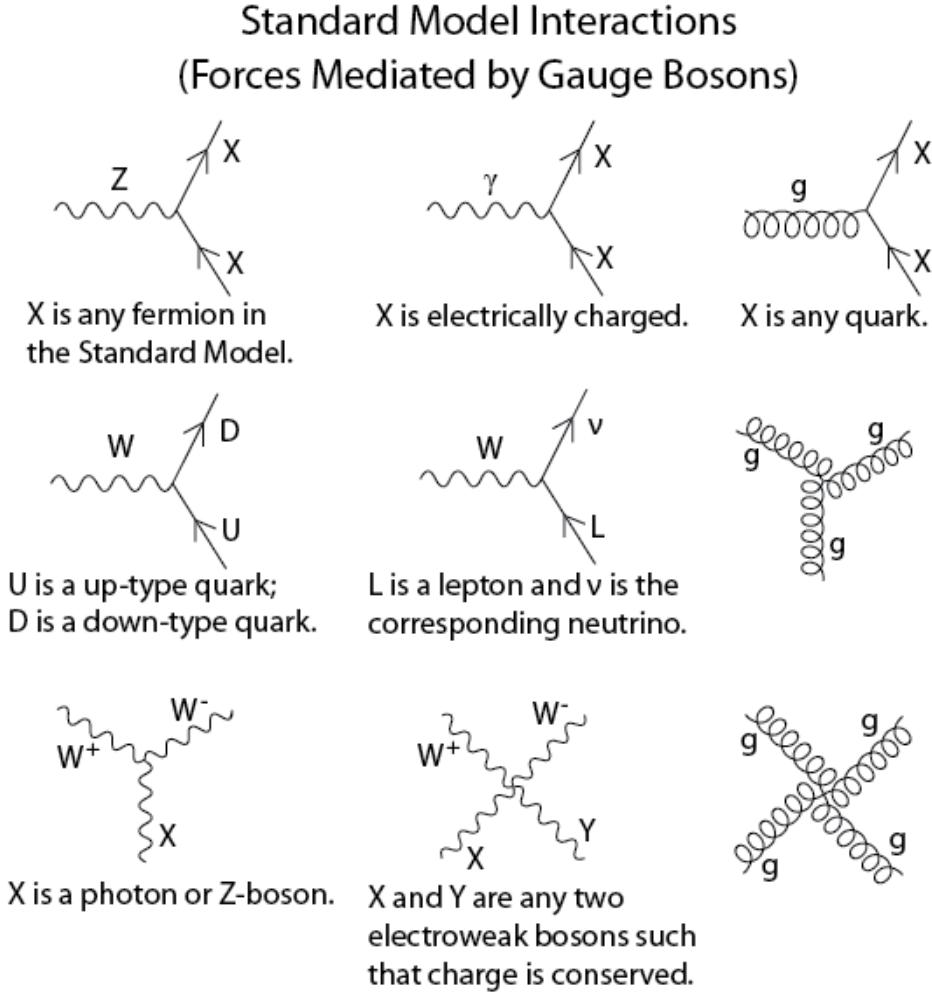
253 The electromagnetic force, mediated by the photon, interacts with via a three-  
254 point coupling all charged particles in the Standard Model. The photon thus interacts  
255 with all the quarks, the charged leptons, and the charged  $W^\pm$  bosons.

256 The weak force is mediated by three particles : the  $W^\pm$  and the  $Z^0$ . The  $Z^0$  can  
257 interact with all fermions via a three-point coupling. A real  $Z_0$  can thus decay to  
258 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



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

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

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

## 274 2.3 Deficiencies of the Standard Model

275 At this point, it is quite easy to simply rest on our laurels. This relatively simple  
276 theory is capable of explaining a very wide range of phenomena, which ultimately  
277 break down only to combinations of nine diagrams shown in Fig.2.3. Unfortunately,  
278 there are some unexplained problems with the Standard Model. We cannot go  
279 through all of the potential issues in this thesis, but we will motivate the primary  
280 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 :

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

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

286 An additional issue is the lack of *gauge coupling unification*. The couplings of  
287 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

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

290 One would hope this behavior was repeated between the electroweak forces and the  
 291 strong force at some suitable energy scale. The Standard Model does automatically  
 292 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.

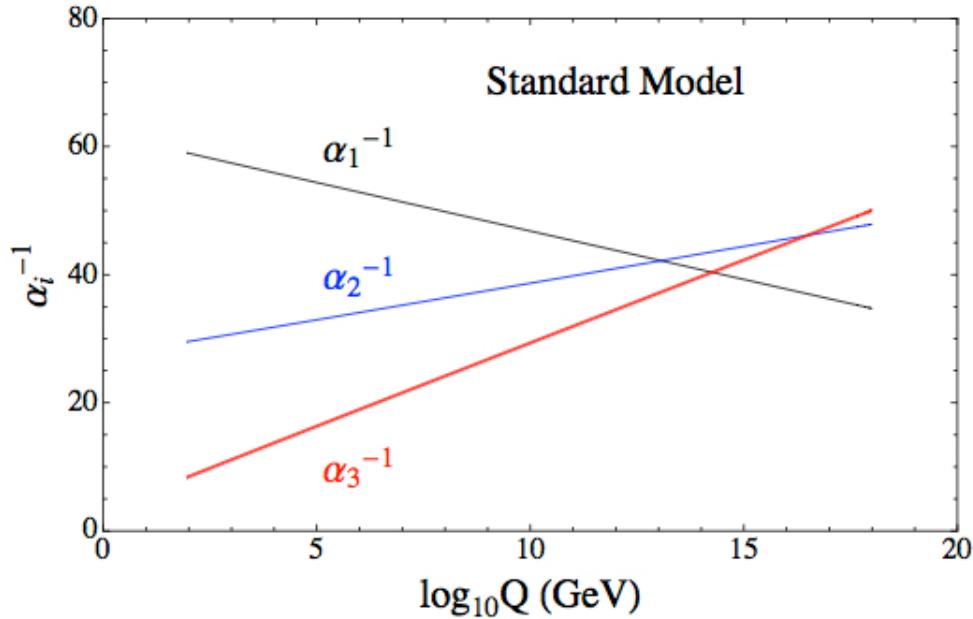
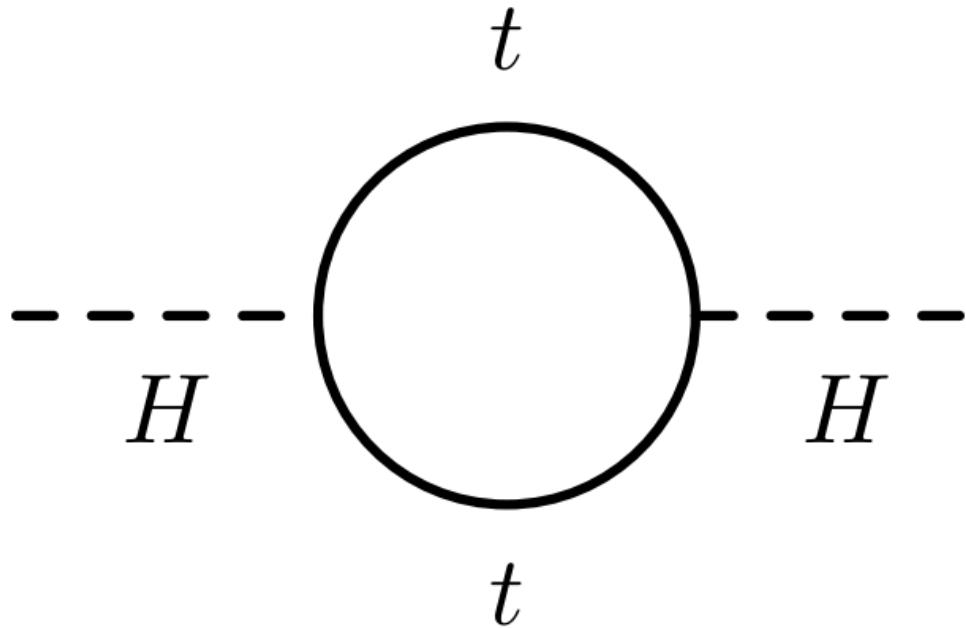


Figure 2.5: The dominant quantum loop correction to the Higgs mass in the Standard Model.



case, we expect the corrections to the Higgs mass like

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

293 To achieve the miraculous cancellation required to get the observed Higgs mass of  
294 125 GeV, one needs to then set the bare Higgs mass  $m_0$ , our input to the Standard  
295 Model Lagrangian, itself to a *precise* value  $\sim 10^{19}$  GeV. This extraordinary level of  
296 parameter finetuning is quite undesirable, and within the framework of the Standard  
297 Model, there is little that can be done to alleviate this issue.

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

309 In the next chapter, we will see how these problems can be alleviated by the theory  
310 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 → $0$
u	c	t	g	H										
up	charm	top	gluon	Higgs boson										
d	s	b	$\gamma$											
down	strange	bottom	photon											
e	$\mu$	$\tau$	Z											
electron	muon	tau	Z boson											
$\nu_e$	$\nu_\mu$	$\nu_\tau$	W											
electron neutrino	muon neutrino	tau neutrino	W boson											



311

## Chapter 3

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312

### *Supersymmetry*

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

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

322 **3.1 Supersymmetric theories : from space to  
323 superspace**

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

325 We begin the theoretical motivation for supersymmetry by citing the “no-go” theorem  
326 of Coleman and Mandula [65]. This theorem forbids *spin-charge unification*; it  
327 states that all quantum field theories which contain nontrivial interactions must be  
328 a direct product of the Poincarégroup of Lorentz symmetries, the internal product  
329 from of gauge symmetries, and the discrete symmetries of parity, charge conjugation,  
330 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].

### 337 Supersymmetry transformations

A *supersymmetric* transformation  $Q$  transforms a bosonic state into a fermionic state, 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] :

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

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

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

340 **Supermultiplets**

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

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

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

366 gauge transformation properties.

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

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

## 376 3.2 Minimally Supersymmetric Standard Model

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

### 383 Chiral supermultiplets

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

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

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

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

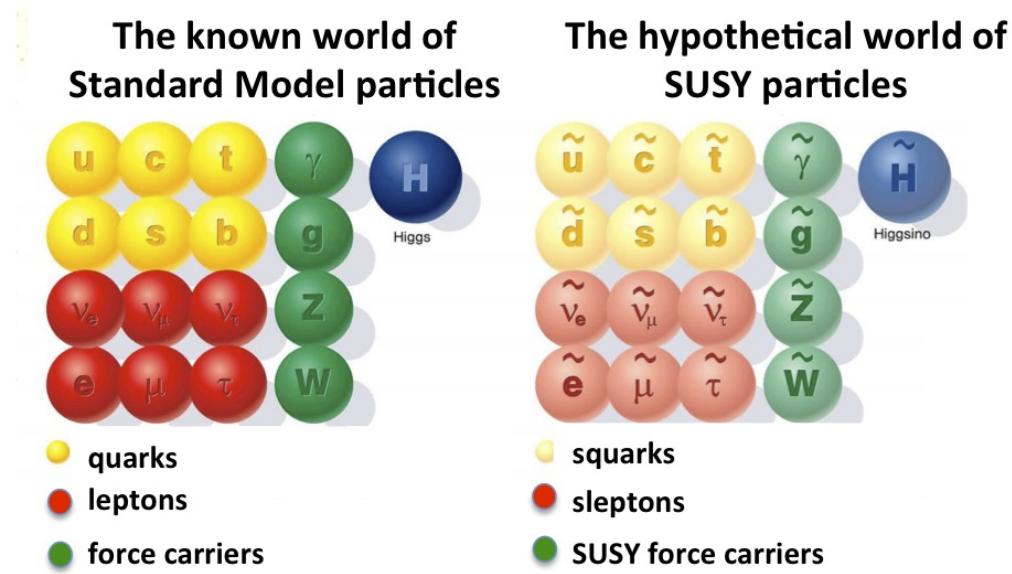
(3.8)

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



### 403 Gauge supermultiplets

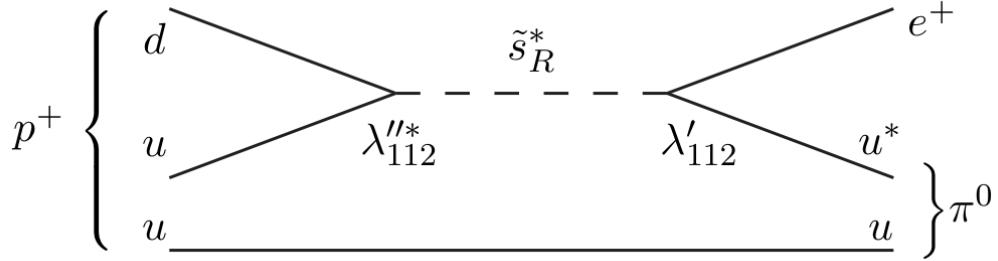
404 The superpartners of the gauge bosons must all be in gauge supermultiplets since  
 405 they contain a spin-1 particle. Collectively, we refer to the superpartners of the  
 406 gauge bosons as the gauginos.

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

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

416 At this point, it's important to take a step back. Where are these particles?  
 417 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.



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

## 421 **$R$ -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

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

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

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

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

430 **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 :

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

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

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

- 435 • Mass terms for the scalar components of the chiral supermultipletss  
 436 • Mass terms for the Weyl spinor components of the gauge supermultipletss  
 437 • 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

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

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

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

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

438 where we have introduced the following notations :

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

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

441 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.

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

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

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

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

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

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

463      

### 3.3 Phenomenology

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

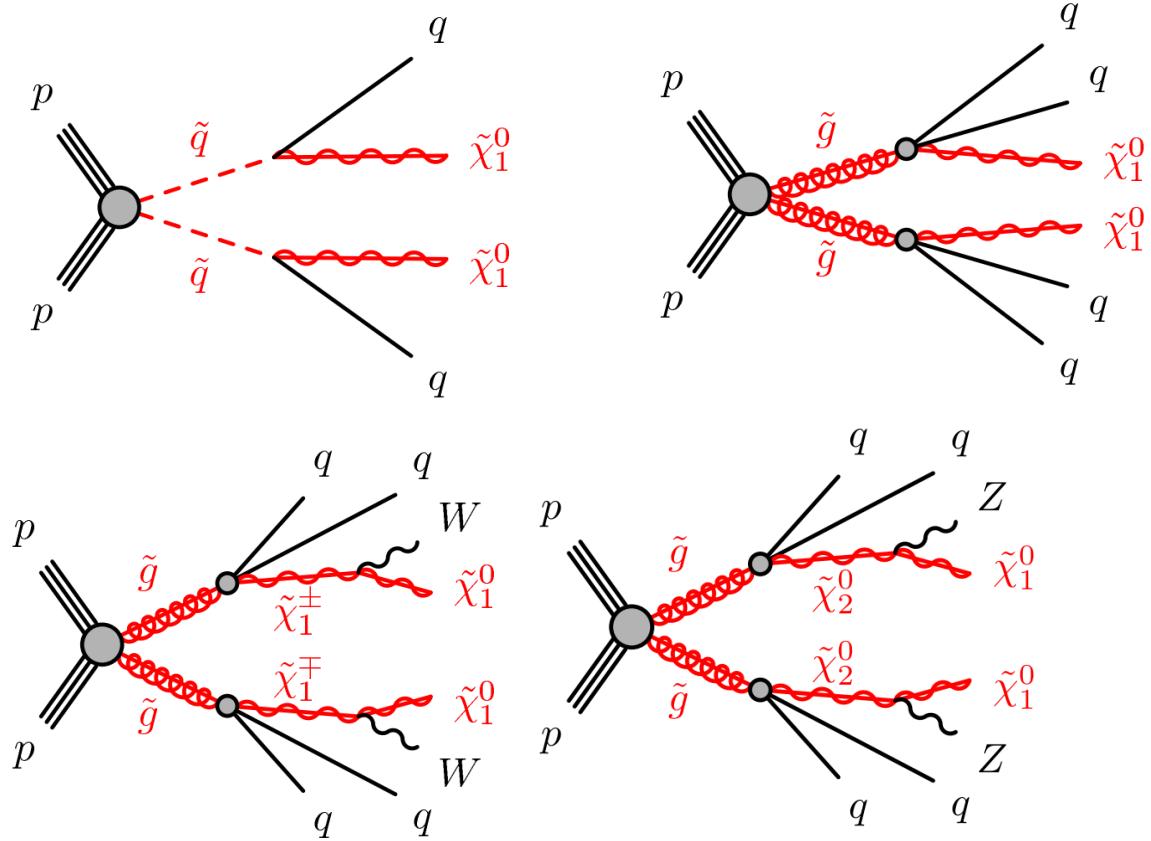
466      As noted above in Sec.3.2, the assumption of *R*–parity has important conse-  
467      quences for MSSM phenomenology. The SM particles have  $R = 1$ , while the sparticles  
468      all have  $R = -1$ . Simply, this is the “charge” of supersymmetry. Since the particles of  
469      LHC collisions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

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



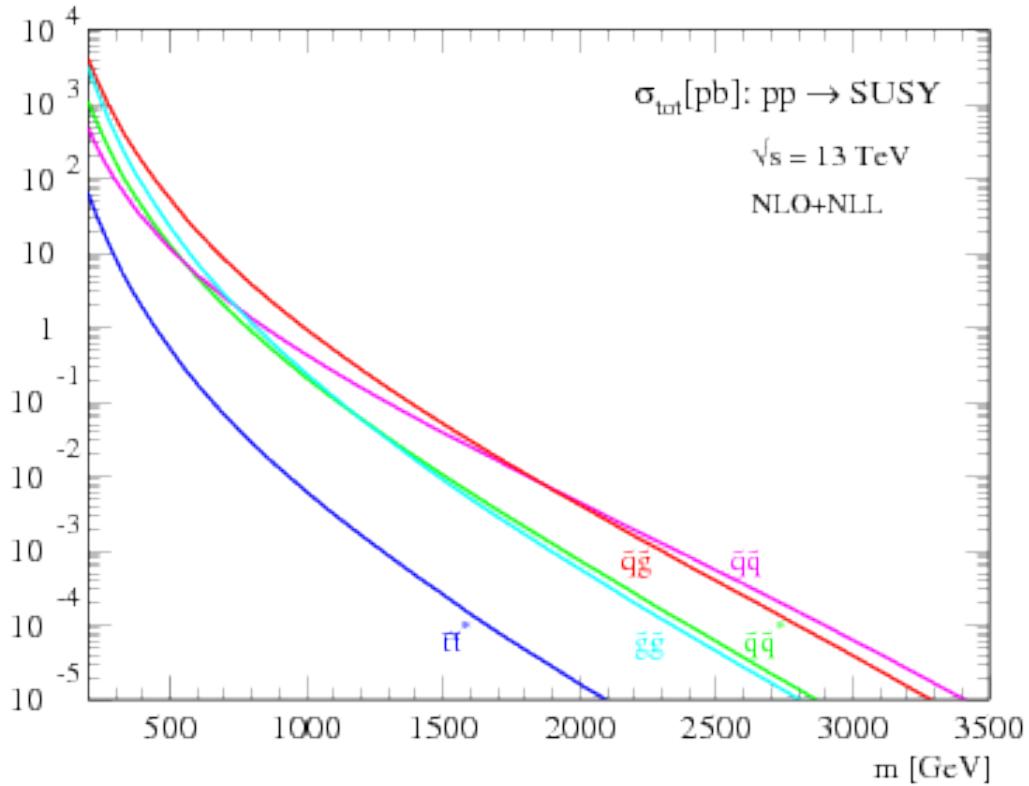
483 of  $SU(3)C$ .

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

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

494 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.

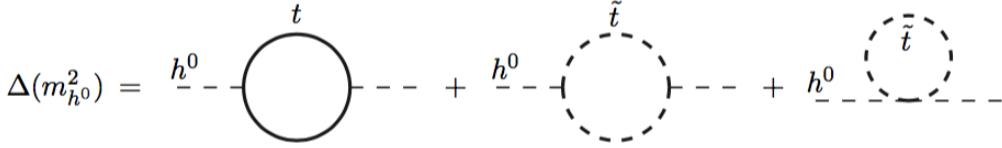


495 *models*. These models make certain assumptions which allow easy comparisons of  
 496 results by theorists and rival experimentalists. In the context of this thesis, the  
 497 simplified models will make assumptions about the branching ratios described in the  
 498 preceding paragraphs. In particular,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

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<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



### 503 3.4 How SUSY solves the problems with the SM

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

#### 506 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

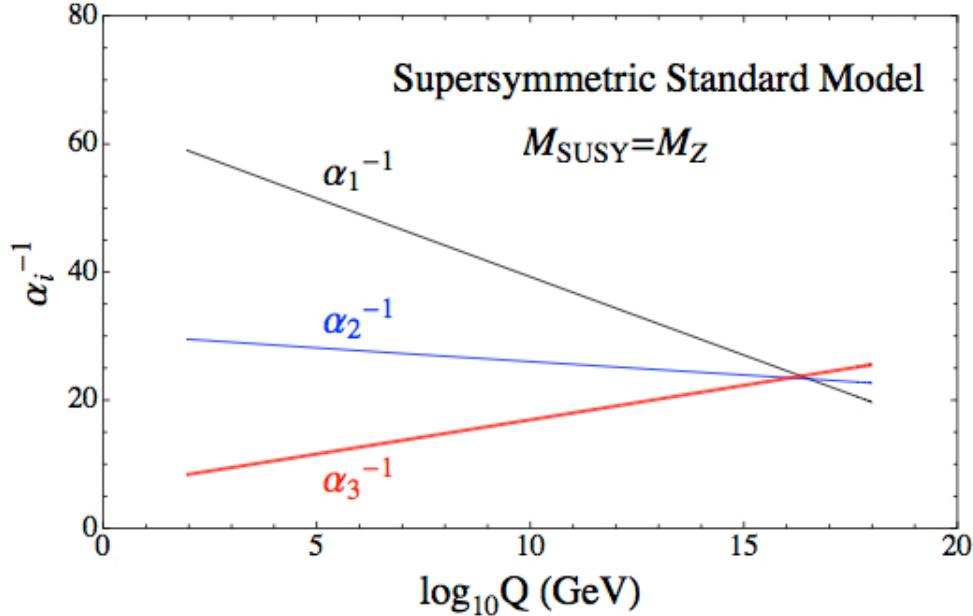
$$\delta m_H^2 \approx \left( \frac{m_t}{8\pi^2 \langle \phi \rangle_{VEV}} \right)^2 \Lambda_{Planck}^2. \quad (3.16)$$

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

#### 512 Gauge coupling unification

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

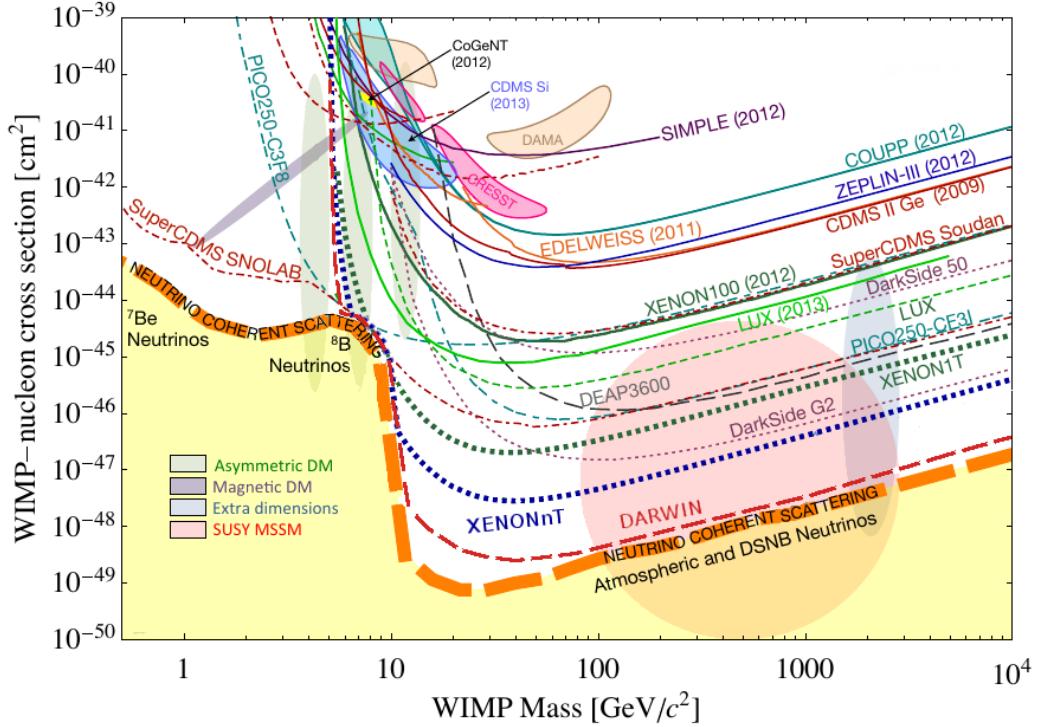
Figure 3.6: The running of Standard Model gauge couplings; compare to Fig.2.4. The MSSM gauge couplings nearly intersect at high energies.



## 518 Dark matter

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

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



## 529 3.5 Conclusions

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



*The Large Hadron Collider*

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

549 **4.1 Basics of Accelerator Physics**

550 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

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

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

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

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

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

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

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

570 The linear relation between  $r$  and  $p$ , or alternatively  $B$  and  $p$  have important  
571 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.

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

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

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

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

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

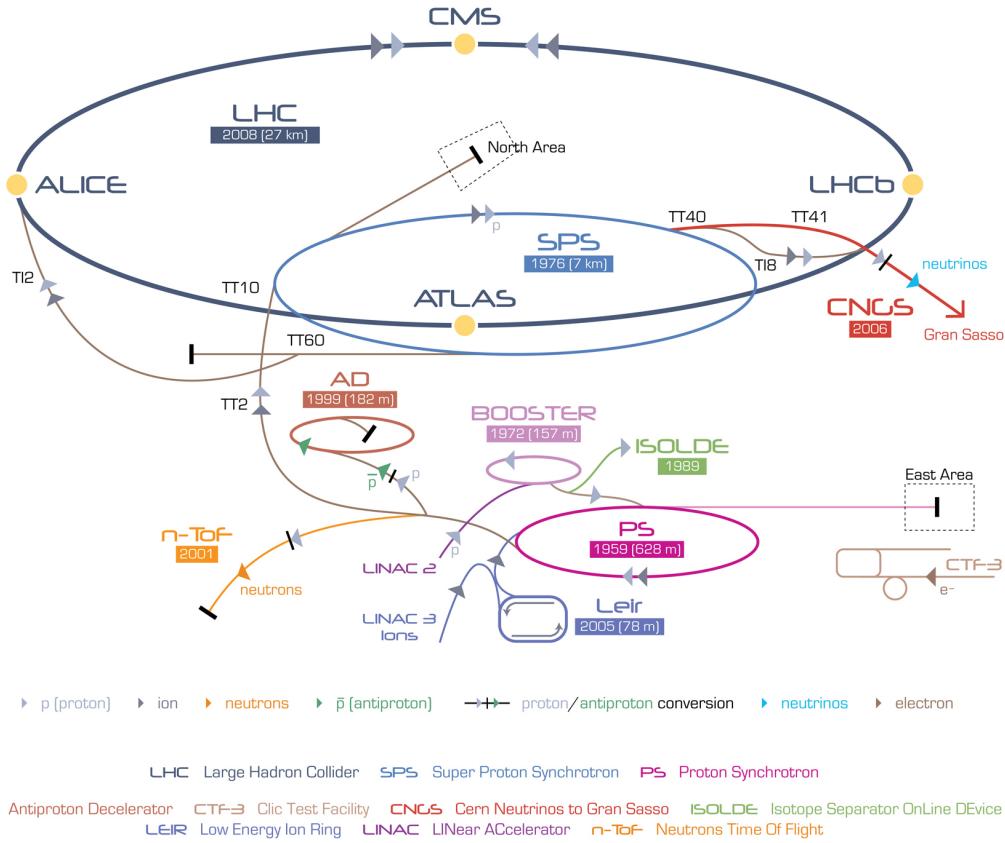
$$N_P = \sigma_P \int L. \quad (4.6)$$

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

## 584 4.2 Accelerator Complex

585 The Large Hadron Collider is the last accelerator in a chain of accelerators which  
586 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.

598 **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 :

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

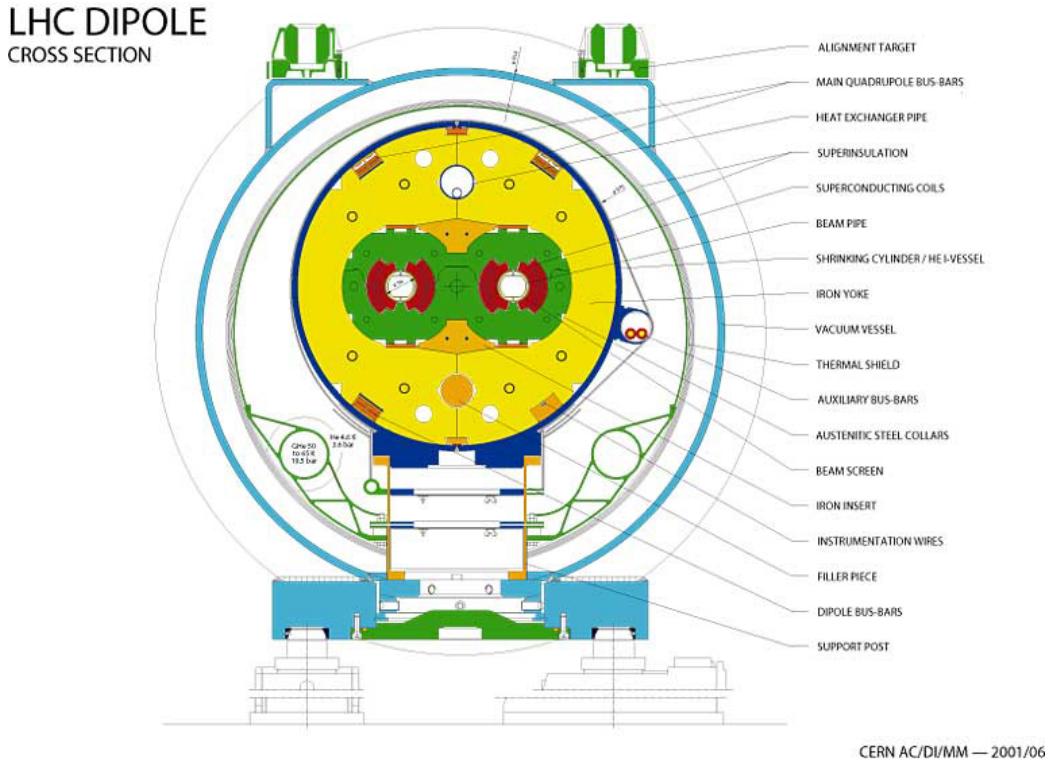
$$\rightarrow B = \frac{p}{rc} = 5 \text{ T} \quad (4.8)$$

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

604 **Magnets**

605 There are many magnets used by the LHC machine, but the most important are the  
606 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.



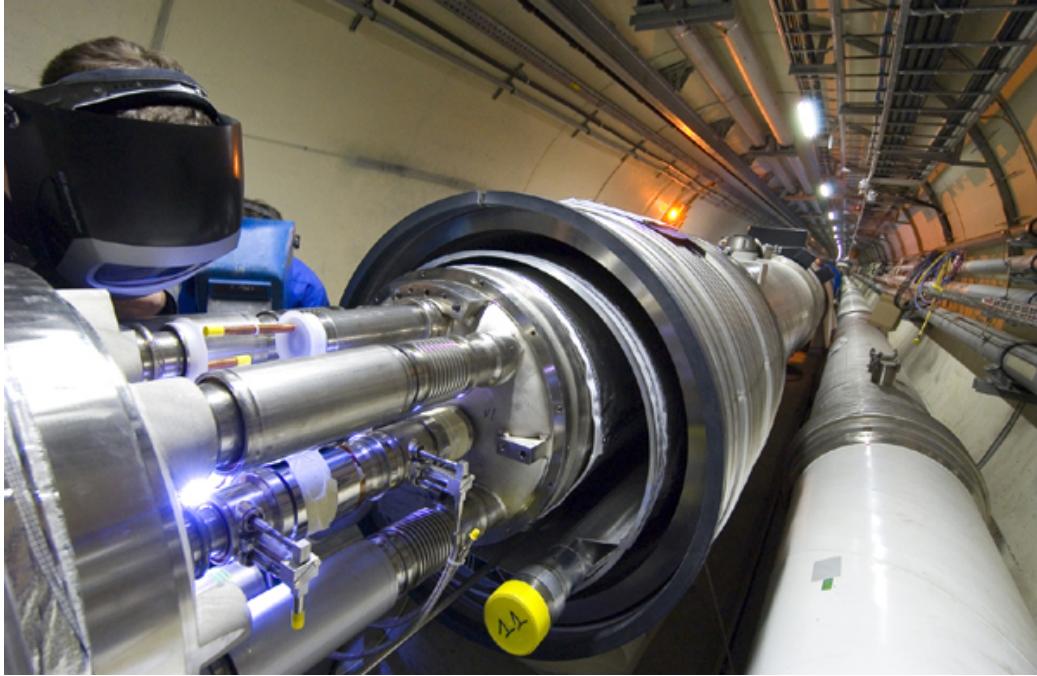
607 Fig.4.3.

608 The magnets are made of Niobium and Titanium. The maximum field strength is  
 609 10 T when cooled to 1.9 Kelvin. The magnets are cooled by superfluid helium, which  
 610 is supplied by a large cryogenic system. Due to heating between the eight helium  
 611 refrigerators and the beampipe, the helium is cooled in the refrigerators to 1.8 K.

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

616 The dipole magnets are 16.5 meters long with a diameter of 0.57 meters. There  
 617 are two individual beam pipes inside each magnet, which allows the dipoles to house  
 618 the beams travelling in both directions around the LHC ring. They curve slightly,  
 619 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.



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

## 622 **4.4 Dataset Delivered by the LHC**

623 In this thesis, we analyze the data delivered by the LHC to ATLAS in the 2015 and  
624 2016 datasets. The beam parameters relevant to this dataset are available in Table  
625 [4.1](#).

626 The peak instantaneous luminosity delivered in 2015 (2016) was  $L =$   
627  $5.2(11) \text{ cm}^{-2}\text{s}^{-1} \times 10^{33}$ . One can note that the instantaneous luminosity delivered in  
628 the 2016 dataset exceeds the design luminosity of the LHC. The total integrated  
629 luminosity delivered was  $13.3 \text{ fb}^{-1}$ . In Figure [4.4](#), we display the integrated luminosity  
630 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.

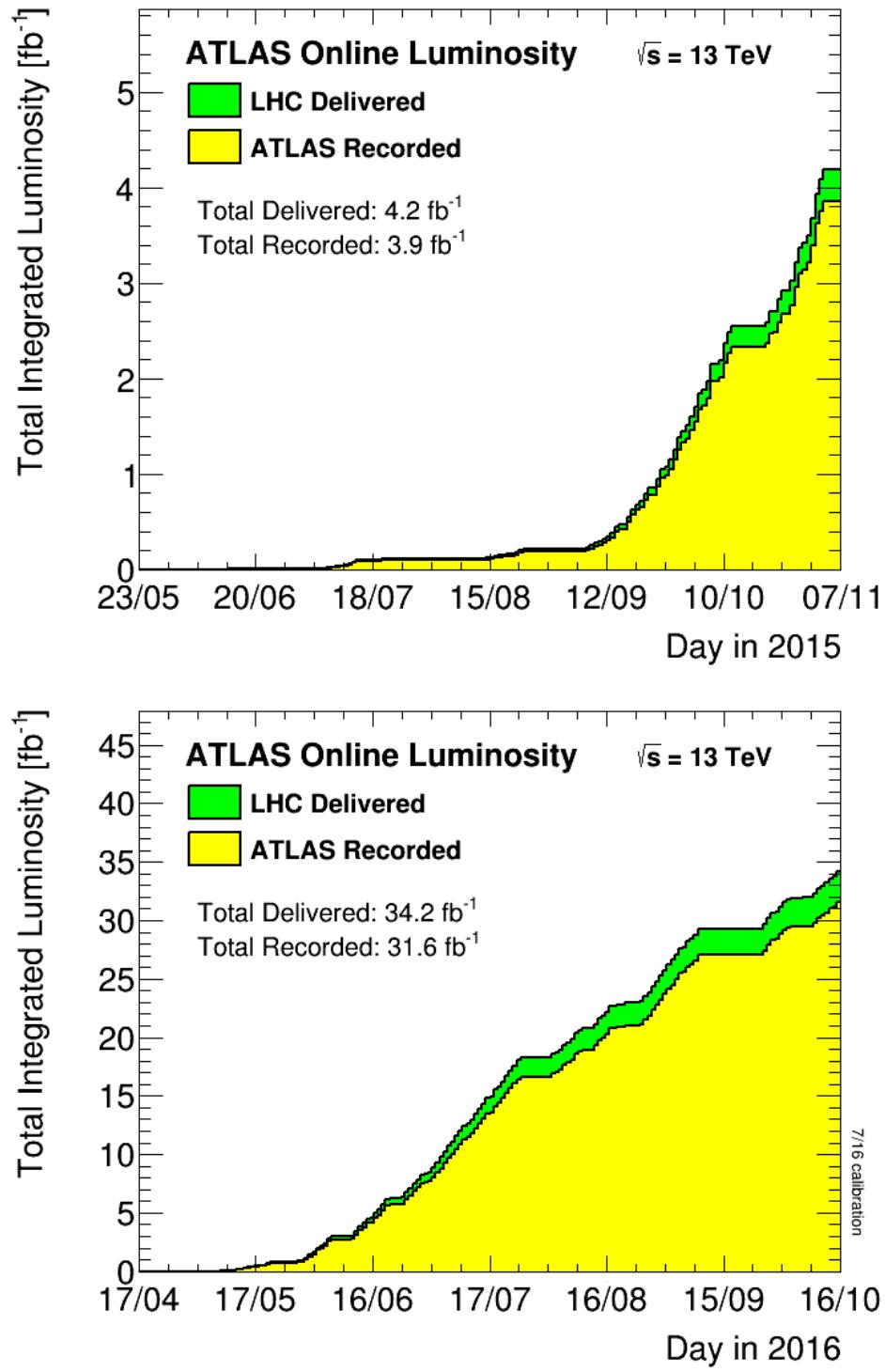
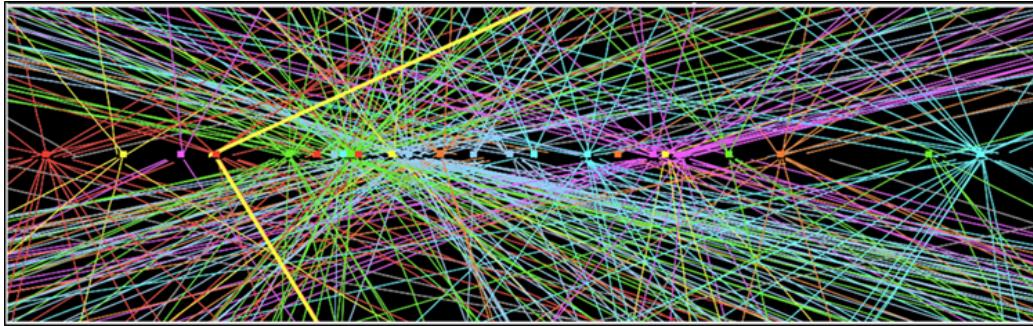


Table 4.1: Beam parameters of the Large Hadron Collider.

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.



## 631 Pileup

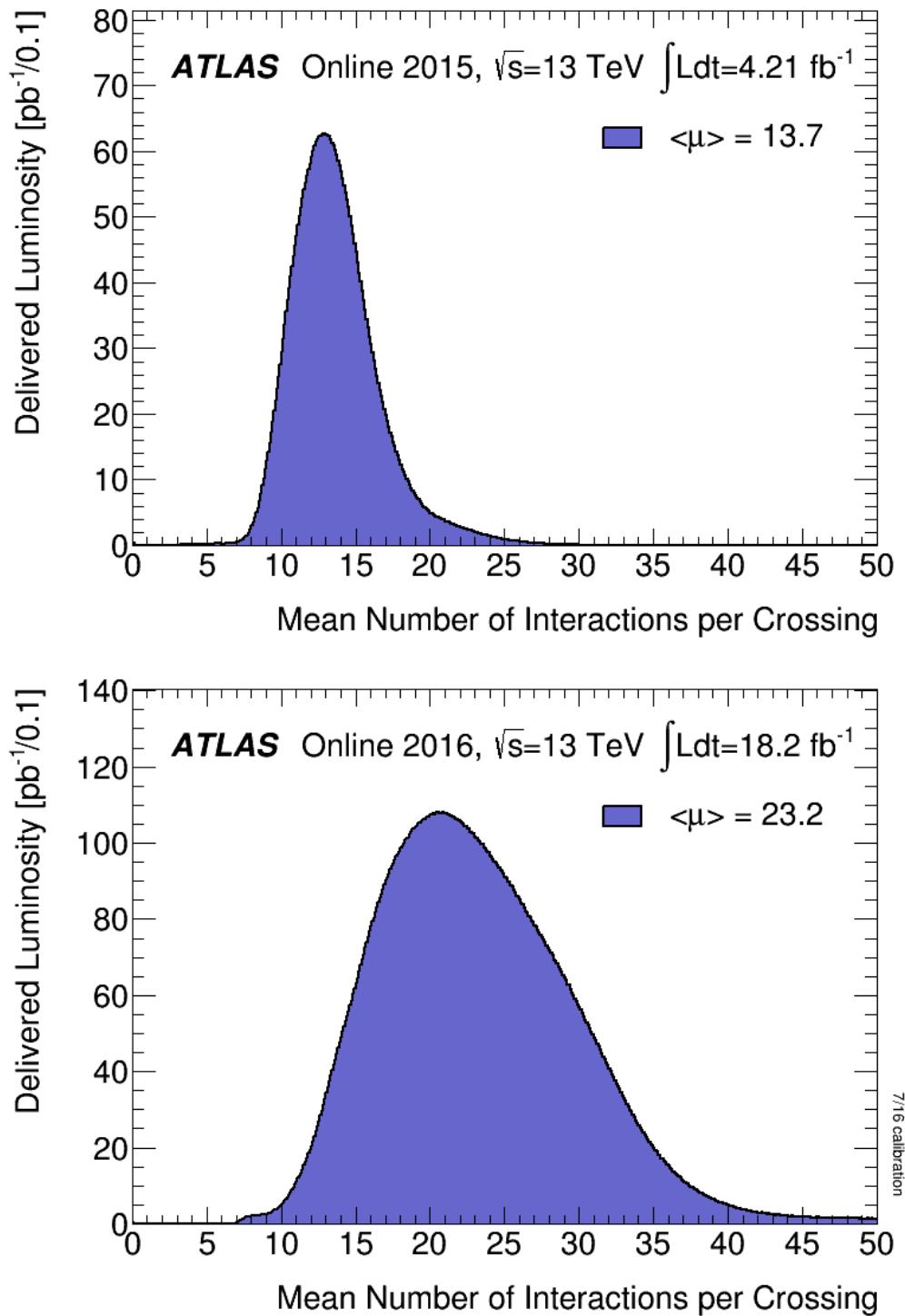
632 *Pileup* is the term for the additional proton-proton interactions which occur during  
 633 each bunch crossing of the LHC. At the beginning of the LHC physics program, there  
 634 had not been a collider which averaged more than a single interaction per bunch  
 635 crossing. In the LHC, each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

643        We quantify in-time pileup by the number of “primary”<sup>2</sup> vertices in a particular  
644    event. To quantify the out-of-time pileup, we use the average number of interactions  
645    per bunch crossing  $\langle \mu \rangle$  over some human-scale time. In Figure 4.6, we show the  
646    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*

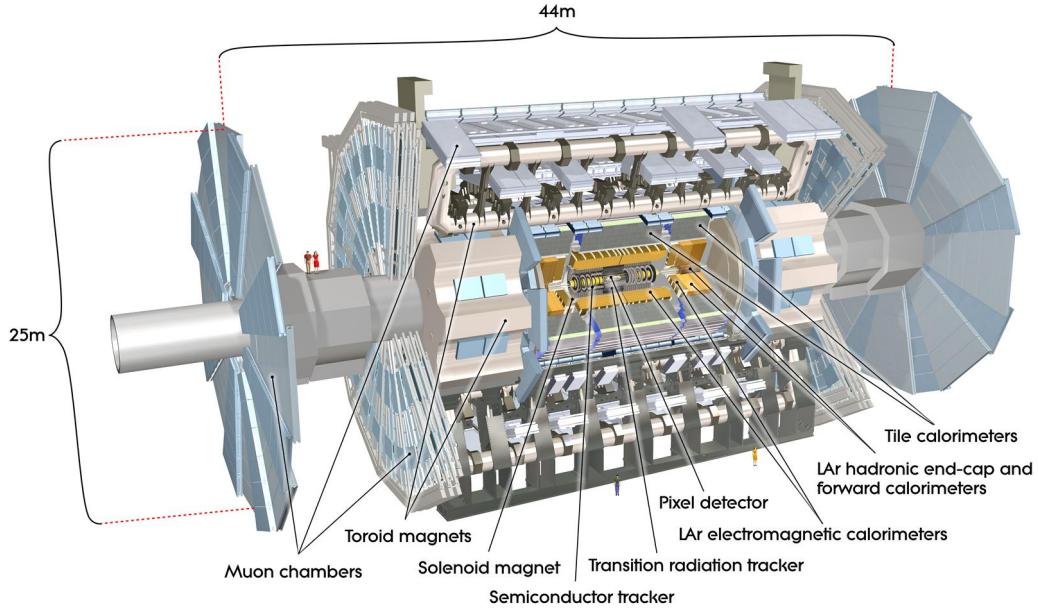
649 The dataset analyzed in this thesis was taken by the ATLAS detector [88], which is  
 650 located at the “Point 1” cavern of the LHC beampipe, just across the street from  
 651 the main CERN campus. The much-maligned acronym stands for *A Toroidal LHC*  
 652 *ApparatuS*. ATLAS is a massive cylindrical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

661 The *inner detector* (ID) lies closest to the collision point, and contains three  
 662 separate subdetectors. It provides pseudorapidity<sup>1</sup>coverage of  $|\eta| < 2.5$  for charged  
 663 particles to interact with the tracking material. The tracks reconstructed from the  
 664 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

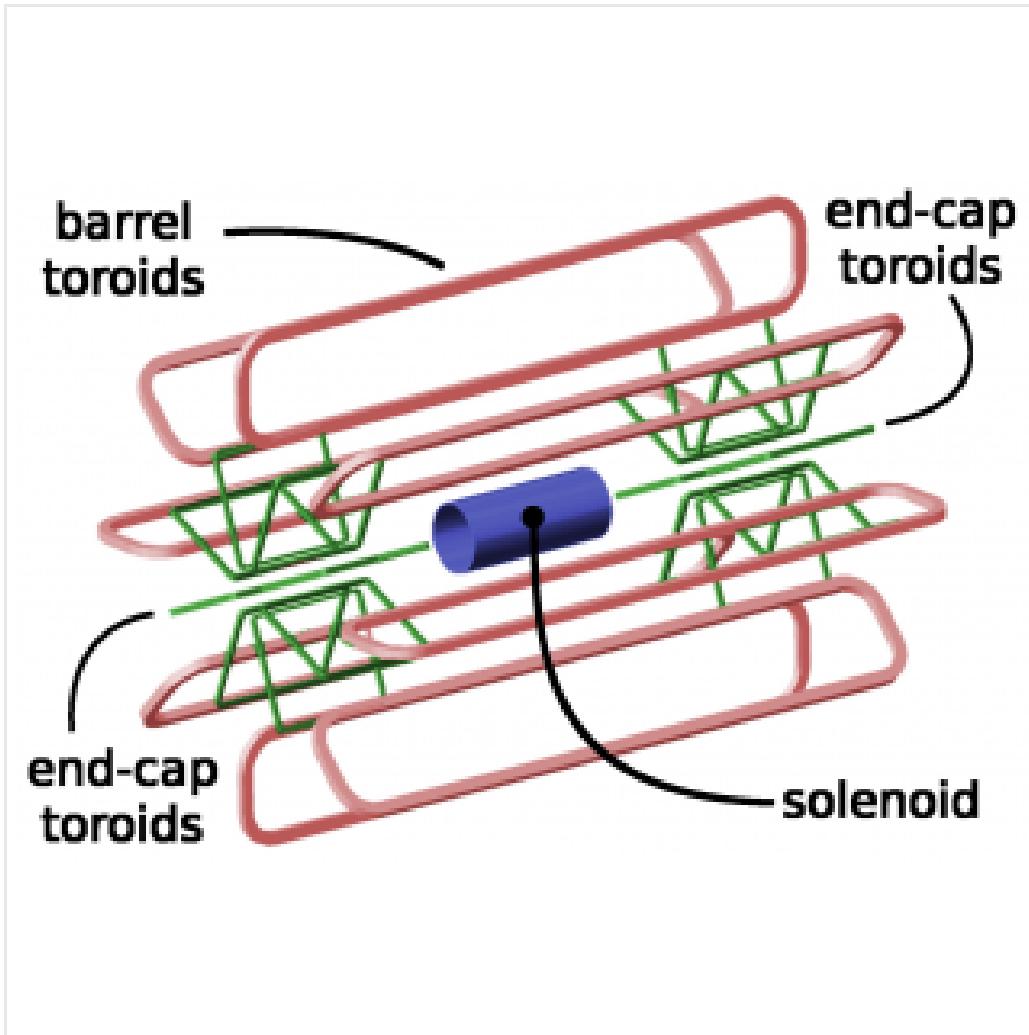


and to determine the momenta of charged particles. The ATLAS *calorimeter* consists of two subdetectors, known as the *electromagnetic* and *hadronic* calorimeters. These detectors stop particles in their detector material, and measure the energy deposition inside, which measures the energy of the particles 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

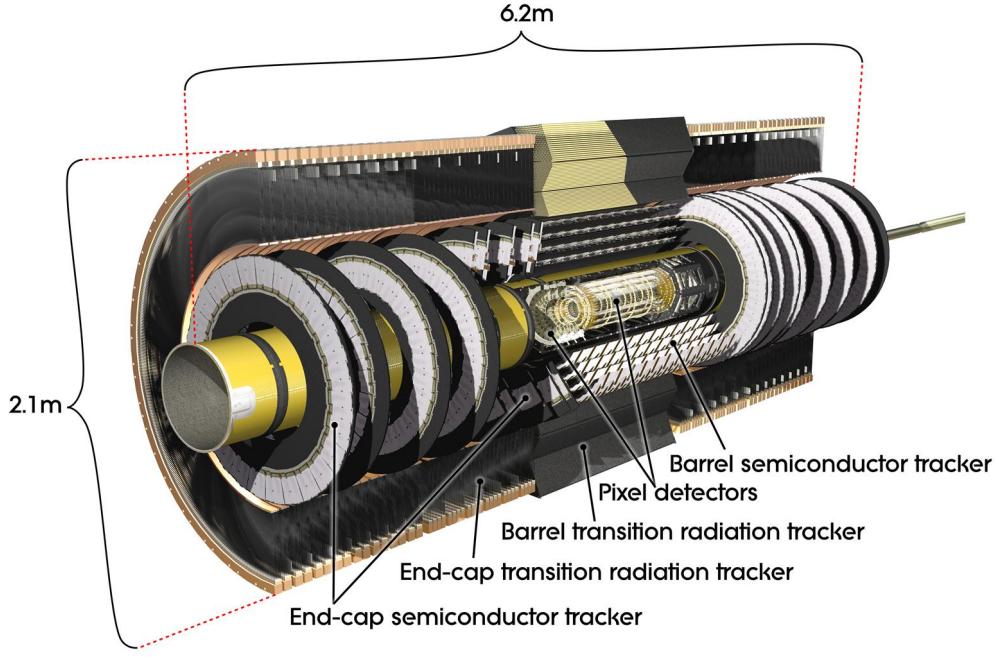
Figure 5.2: The ATLAS magnet system



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

680 The ATLAS central solenoid [89] is a 2.3 m diameter, 5.3 m long solenoid at the  
681 center of the ATLAS detector. It produces a uniform magnetic field of 2 T; this  
682 strong field is necessary to accurately measure the charged particles in this field.  
683 An important design constraint for the central solenoid was the decision to place  
684 it in between the inner detector and the calorimeters. To avoid excessive impacts  
685 on measurements in the calorimetry, the central solenoid must be as transparent as

Figure 5.3: The ATLAS inner detector



686 possible<sup>2</sup>.

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

## 692 5.2 Inner Detector

693 The ATLAS inner detector consists of three separate tracking detectors, which are  
694 known as, in order of increasing distance from the interaction point, the Pixel Detector  
cite 695 [90], Semiconductor Tracker (SCT) , and the Transition Radiation Tracker (TRT) .

cite 696 When charged particles pass through these tracking layers, they produce *hits*, which

---

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

697 using the known 2 T magnetic field, allows the reconstruction of *tracks*. Tracks  
698 are used as inputs for reconstruction of many higher-level physics objects, such as  
699 electrons, muons, photons, and  $E_T^{\text{miss}}$ . Accurate track reconstruction is thus crucial  
700 for precise measurements of charged particles.

701 **Pixel Detector**

702

schematic

703 The ATLAS pixel detector consists four layers of silicon “pixels”. This refers to the  
704 segmentation of the active medium into the pixels; compare to the succeeding silicon  
705 detectors, which will use silicon “strips”. This provides precise 3D hit locations. The  
706 layers are known as the “Insertable”<sup>3</sup>B-Layer (IBL), the B-Layer (or Layer-0), Layer-  
707 1, and Layer-2, in order of increasing distance from the interaction point. These  
708 layers are very close to the interaction point, and therefore experience a large amount  
709 of radiation.

710 Layer-1, Layer-2, and Layer-3 were installed with the initial construction of  
711 ATLAS. They contain front-end integrated electronics (FEI3s) bump-bonded to 1744  
712 silicon modules; each module is 250  $\mu\text{m}$  in thickness and contains 47232 pixels. These  
713 pixels have planar sizes of 50 x 400  $\mu\text{m}^2$  or 50 x 600  $\mu\text{m}^2$ , to provide highly accurate  
714 location information. The FEI3s are mounted on long rectangular structures known  
715 as staves, which encircle the beam pipe. A small tilt to each stave allows full coverage  
716 in  $\phi$  even with readout systems which are installed. These layers are at radia of 50.5  
717 mm, 88.5 mm, and 122.5 mm from the interaction point.

718 The IBL was added to ATLAS after Run1 in 2012 at a radius of 33 mm from the  
719 interaction point. The entire pixel detector was removed from the center of ATLAS  
720 to allow an additional pixel layer to be installed. The IBL was required to preserve

---

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

721 the integrity of the pixel detector as radiation damage leads to inoperative pixels in  
722 the other layers. The IBL consists of 448 FEI4 chips, arranged onto 14 staves. Each  
723 FEI4 has 26880 pixels, of planar size  $50 \times 250 \mu\text{m}$ . This smaller granularity was  
724 required due to the smaller distance to the interaction point.

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

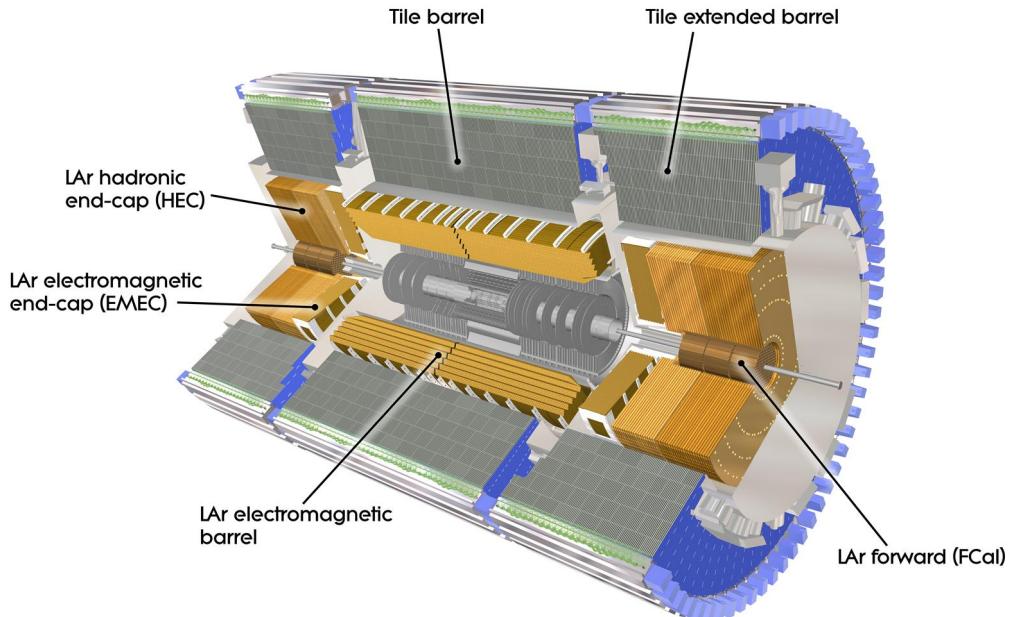
## 727 Semiconductor Tracker

schematic 728 The SCT is directly beyond Layer-2 of the pixel detector. This is a silicon strip  
729 detector, which do not provide the full 3D information of the pixel detector. The  
730 dual-sensors of the SCT contain  $2 \times 768$  individual strips; each strip has area  $6.4$   
731  $\text{cm}^2$ . The SCT dual-sensor is then double-layered, at a relative angle of 40 mrad;  
732 together these layers provide the necessary 3D information for track reconstruction.  
733 There are four of these double-layers, at radia of 284 mm, 355 mm, 427 mm, and 498  
734 mm. These double-layers provide hits comparable to those of the pixel detector, and  
735 we have four additional hits to reconstruct tracks for each charged particle.

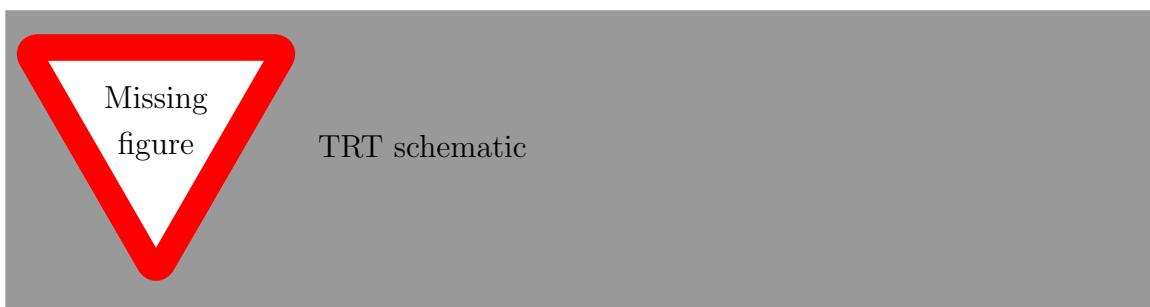
## 736 Transition Radiation Tracker

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

Figure 5.4: The ATLAS calorimeter



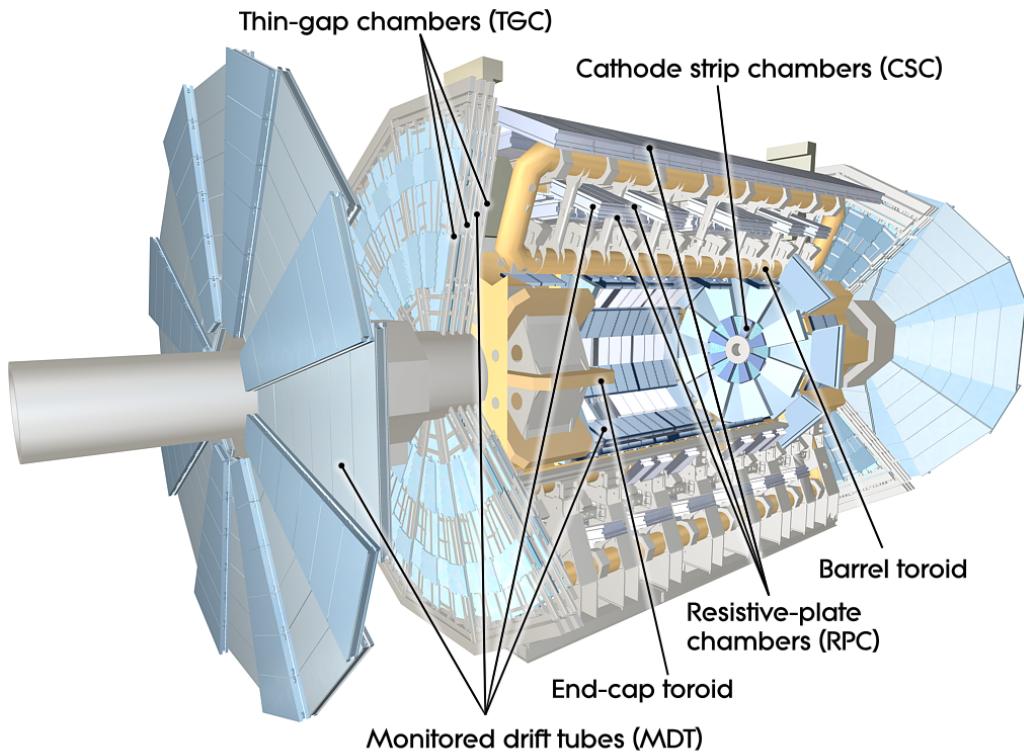
745      The TRT is so-named due to the *transition radiation* (TR) it induces. Due to  
746      the dielectric difference between the gas and tubes, TR is induced. This is important  
747      for distinguishing electrons from their predominant background of minimum ionizing  
748      particles. Generally,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



751

752

Figure 5.5: The ATLAS muon spectrometer



<sup>753</sup> **5.3 Calorimeter**

<sup>754</sup> **Electromagnetic Calorimeter**

<sup>755</sup> **Hadronic Calorimeter**

<sup>756</sup> **5.4 Muon Spectrometer**

<sup>757</sup> **5.5 Trigger System**

758

## Chapter 6

---

759

### *The Recursive Jigsaw Technique*

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

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

### **763 6.1 Razor variables**

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

### **767 6.2 SuperRazor variables**

### **768 6.3 The Recursive Jigsaw Technique**

### **769 6.4 Variables used in the search for zero lepton**

770 **SUSY**



771

## Chapter 7

---

772

### *Title of Chapter 1*



773

## Chapter 8

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774

### *Title of Chapter 1*

775 Here you can write some introductory remarks about your chapter. I like to give each

776 sentence its own line.

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

## **778 8.1 Object reconstruction**

### **779 Photons, Muons, and Electrons**

### **780 Jets**

### **781 Missing transverse momentum**

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

783 **8.2 Signal regions**

784 **Gluino signal regions**

785 **Squark signal regions**

786 **Compressed signal regions**

787 **8.3 Background estimation**

788 **Z vv**

789 **W ev**

790 **ttbar**

*Title of Chapter 1*

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

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

796 **9.1 Statistical Analysis**

797 maybe to be moved to an appendix

798 **9.2 Signal Region distributions**

799 **9.3 Pull Plots**

800 **9.4 Systematic Uncertainties**

801 **9.5 Exclusion plots**



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802

## *Conclusion*

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

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

806 **9.6 New Section**

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



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810

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1047

## *The Standard Model*

1048 In this appendix, we provide a brief overview of the basic ingredients involved in  
1049 construction of the Standard Model Lagrangian : quantum field theory, symmetries,  
1050 and symmetry breaking.

## **1051 Quantum Field Theory**

1052

1053 In this section, we provide a brief overview of the necessary concepts from  
1054 Quantum Field Theory (QFT).

cite Yuval's  
lectures  
and notes  
somehow

1055 In modern physics, the laws of nature are described by the “action”  $S$ , with the  
1056 imposition of the principle of minimum action. The action is the integral over the cite  
1057 spacetime coordinates of the “Lagrangian density”  $\mathcal{L}$ , or Lagrangian for short. The  
1058 Lagrangian is a function of “fields”; general fields will be called  $\phi(x^\mu)$ , where the  
1059 indices  $\mu$  run over the space-time coordinates. We can then write the action  $S$  as

$$S = \int d^4x \mathcal{L}[\phi_i(x^\mu), \partial_\mu \phi_i(x^\mu)] \quad (9.1)$$

1060 where we have an additional summation over  $i$  (of the different fields). Generally,  
1061 we impose the following constraints on the Lagrangian :

- 1062 1. Translational invariance - The Lagrangian is only a function of the fields  $\phi$  and  
1063 their derivatives  $\partial_\mu \phi$
- 1064 2. Locality - The Lagrangian is only a function of one point  $x_\mu$  in spacetime.

1065        3. Reality condition - The Lagrangian is real to conserve probability.

1066        4. Lorentz invariance - The Lagrangian is invariant under the Poincarégroup of  
1067              spacetime.

1068        5. Analyticity - The Lagrangian is an analytical function of the fields; this is to  
1069              allow the use of perturbation theory.

1070        6. Invariance and Naturalness - The Lagrangian is invariant under some internal  
1071              symmetry groups; in fact, the Lagrangian will have *all* terms allowed by the  
1072              imposed symmetry groups.

1073        7. Renormalizability - The Lagrangian will be renormalizable - in practice, this  
1074              means there will not be terms with more than power 4 in the fields.

1075        The key item from the point of view of this thesis is that of “Invariance and  
1076              Natural”. We impose a set of “symmetries” and then our Lagrangian is the most  
1077              general which is allowed by those symmetries.

## 1078        Symmetries

1079        Symmetries can be seen as the fundamental guiding concept of modern physics.  
1080        Symmetries are described by “groups”. To illustrate the importance of symmetries  
1081              and their mathematical description, groups, we start here with two of the simplest  
1082              and most useful examples :  $\mathbb{Z}_2$  and  $U(1)$ .

### 1083        $\mathbb{Z}_2$ symmetry

1084         $\mathbb{Z}_2$  symmetry is the simplest example of a “discrete” symmetry. Consider the most  
1085              general Lagrangian of a single real scalar field  $\phi(x_\mu)$

$$\mathcal{L}_\phi = \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - \frac{m^2}{2}\phi^2 - \frac{\mu}{2\sqrt{2}}\phi^3 - \lambda\phi^4 \quad (9.2)$$

Now we *impose* the symmetry

$$\mathcal{L}(\phi) = \mathcal{L}(-\phi) \quad (9.3)$$

1086 This has the effect of restricting the allowed terms of the Lagrangian. In particular,  
 1087 we can see the term  $\phi^3 \rightarrow -\phi^3$  under the symmetry transformation, and thus must  
 1088 be disallowed by this symmetry. This means under the imposition of this particular  
 1089 symmetry, our Lagrangian should be rewritten as

$$\mathcal{L}_\phi = \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - \frac{m^2}{2}\phi^2 - \lambda\phi^4 \quad (9.4)$$

1090 The effect of this symmetry is that the total number of  $\phi$  particles can only change  
 1091 by even numbers, since the only interaction term  $\lambda\phi^4$  is an even power of the field.  
 1092 This symmetry is often imposed in supersymmetric theories, as we will see in Chapter  
 1093 3.

## 1094 **$U(1)$ symmetry**

1095  $U(1)$  is the simplest example of a continuous (or *Lie*) group. Now consider a theory  
 1096 with a single complex scalar field  $\phi = \text{Re } \phi + i \text{Im } \phi$

$$\mathcal{L}_\phi = \delta_{i,j} \frac{1}{2}\partial_\mu\phi_i\partial^\mu\phi_j - \frac{m^2}{2}\phi_i\phi_j - \frac{\mu}{2\sqrt{2}}\phi_i\phi_j\phi_k\phi_l - \lambda\phi_i\phi_j\phi_k\phi_l \quad (9.5)$$

1097 where  $i, j, k, l = \text{Re}, \text{Im}$ . In this case, we impose the following  $U(1)$  symmetry  
 1098 :  $\phi \rightarrow e^{i\theta}, \phi^* \rightarrow e^{-i\theta}$ . We see immediately that this again disallows the third-order  
 1099 terms, and we can write a theory of a complex scalar field with  $U(1)$  symmetry as

$$\mathcal{L}_\phi = \partial_\mu\phi\partial^\mu\phi^* - \frac{m^2}{2}\phi\phi^* - \lambda(\phi\phi^*)^2 \quad (9.6)$$

## 1100 Local symmetries

1101 The two examples considered above are “global” symmetries in the sense that the  
1102 symmetry transformation does not depend on the spacetime coordinate  $x_\mu$ . We know  
1103 to look at local symmetries; in this case, for example with a local  $U(1)$  symmetry, the  
1104 transformation has the form  $\phi(x_\mu) \rightarrow e^{i\theta(x_\mu)}\phi(x_\mu)$ . These symmetries are also known  
1105 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 9.2 are *not* invariant under a local symmetry transformation

$$\partial_\mu \phi(x_\mu) \rightarrow \partial_\mu(e^{i\theta(x_\mu)}\phi(x_\mu)) = (1 + i\theta(x_\mu))e^{i\theta(x_\mu)}\phi(x_\mu) \quad (9.7)$$

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1106 This leads us to note that the kinetic terms of the Lagrangian are also not invariant  
1107 under a gauge symmetry. This would lead to a model with no dynamics, which is  
1108 clearly unsatisfactory.

1109 Let us take inspiration from the case of global symmetries. We need to define a  
1110 so-called “covariant” derivative  $D^\mu$  such that

$$D^\mu \phi \rightarrow e^{iq\theta(x^\mu)D^\mu}\phi \quad (9.8)$$

$$D^\mu \phi^* \rightarrow e^{-iq\theta(x^\mu)D^\mu}\phi^* \quad (9.9)$$

$$(9.10)$$

1111 Since  $\phi$  and  $\phi^*$  transforms with the opposite phase, this will lead the invariance  
1112 of the Lagrangian under our local gauge transformation. This  $D^\mu$  is of the following  
1113 form

$$D^\mu = \partial_\mu - igqA^\mu \quad (9.11)$$

1115 where  $A^\mu$  is a vector field we introduce with the transformation law

$$A^\mu \rightarrow A^\mu - \frac{1}{g} \partial_\mu \theta \quad (9.12)$$

1116 and  $g$  is the coupling constant associated to vector field. This vector field  $A^\mu$  is  
1117 also known as a “gauge” field.

1118 Since we need to add all allowed terms to the Lagrangian, we define

$$F^{\mu\nu} = A^\mu A^\nu - A^\nu A^\mu \quad (9.13)$$

1119 and then we must also add the kinetic term :

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4} F^{\mu\nu} F_{\mu\nu} \quad (9.14)$$

1120 The most general renormalizable Lagrangian with fermion and scalar fields can  
1121 be written in the following form

$$\mathcal{L} = \mathcal{L}_{\text{kin}} + \mathcal{L}_\phi + \mathcal{L}_\psi + \mathcal{L}_{\text{Yukawa}} \quad (9.15)$$

## 1122 Symmetry breaking and the Higgs mechanism

1123 Here we view some examples of symmetry breaking. We investigate breaking of a  
1124 global  $U(1)$  symmetry and a local  $U(1)$  symmetry. The SM will break the electroweak  
1125 symmetry  $SU(2)xU(1)$ , and in Chapter 3 we will see how supersymmetry must also  
1126 be broken.

1127 There are two ideas of symmetry breaking

- 1128 • Explicit symmetry breaking by a small parameter - in this case, we have a small  
1129 parameter which breaks an “approximate” symmetry of our Lagrangian. An  
1130 example would be the theory of the single scalar field 9.2, when  $\mu \ll m^2$  and

1131       $\mu \ll \lambda$ . In this case, we can often ignore the small term when considering  
 1132      low-energy processes.

1133      • Spontaneous symmetry breaking (SSB) - spontaneous symmetry breaking  
 1134      occurs when the Lagrangian is symmetric with respect to a given symmetry  
 1135      transformation, but the ground state of the theory is *not* symmetric with respect  
 1136      to that transformation. This can have some fascinating consequences, as we  
 1137      will see in the following examples

1138      Symmetry breaking a

1139      **U(1) global symmetry breaking**

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

$$\phi \rightarrow e^{i\theta} \phi \quad (9.16)$$

The Lagrangian for this theory is

$$\mathcal{L} = \partial^\mu \phi^\dagger \partial_\mu \phi + \frac{\mu^2}{2} \phi^\dagger \phi + \frac{\lambda}{4} (\phi^\dagger \phi)^2 \quad (9.17)$$

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

$$\mathcal{L} = \partial^\mu h \partial_\mu h + \partial^\mu \xi d\mu \xi - \frac{\mu^2}{2} (h^2 + \xi^2) - \frac{\lambda}{4} (h^2 + \xi^2)^2 \quad (9.18)$$

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 :

$$V(\phi) = \lambda(\phi^\dagger \phi - v^2/2)^2 \quad (9.19)$$

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

$$2 < \phi^\dagger \phi > = < h^2 + \xi^2 > = v^2 \quad (9.20)$$

1140        We now reach the “breaking” point of this procedure. In the  $(h, \xi)$  plane, the  
 1141   minima form a circle of radius  $v$ . We are free to choose any of these minima to expand  
 1142   our Lagrangian around; the physics is not affected by this choice. For convenience,  
 1143   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

$$\mathcal{L} = \frac{1}{2}\partial_\mu h'\partial^\mu h' + \frac{1}{2}\partial_\mu \xi'\partial^\mu \xi' - \lambda v^2 h'^2 - \lambda v h'(h'^2 + \xi'^2) - \lambda(h'^2 + \xi'^2)^2 \quad (9.21)$$