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A search for sparticles in zero lepton final states

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Submitted in partial fulfillment of the

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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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Acknowledgements

Introduction

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

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

Despite its impressive range of described phenomena, the Standard Model has some theoretical and experimental deficiencies. The SM contains 26 free parameters¹. It would be more theoretically pleasing to understand these free parameters in terms of a more fundamental theory. The major theoretical concern of the Standard Model, as it pertains to this thesis, is the *hierachy problem*[11–15]. The light mass

¹This 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 α_{force}).

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

89 Both of these major issues, as well as numerous others, can be solved by the
90 introduction of *supersymmetry* (SUSY) [15, 23–33]. In supersymmetric theories, each
91 SM particles has a so-called *superpartner*, or sparticle partner, differing from given SM
92 particle by 1/2 in spin. These theories solve the hierachy problem, since the quantum
93 corrections induced from the superpartners exactly cancel those induced by the SM
94 particles. In addition, these theories are usually constructed assuming R -parity,
95 which can be thought of as the “charge” of supersymmetry, with SM particles having
96 $R = 1$ and sparticles having $R = -1$. In collider experiments, since the incoming
97 SM particles have total $R = 1$, the resulting sparticles are produced in pairs. This
98 produces a rich phenomenology, which is characterized by significant hadronic activity
99 and large missing transverse energy (E_T^{miss}), which provide significant discrimination
100 against SM backgrounds [34].

101 Despite the power of searches for supersymmetry where E_T^{miss} is a primary
102 discriminating variable, there has been significant interest in the use of other variables
103 to discriminate against SM backgrounds. These include searches employing variables
104 such as αT , $M_{T,2}$, and the razor variables (M_R, R^2) [35–45]. In this thesis, we
105 will present the first search for supersymmetry using the novel Recursive Jigsaw
106 Reconstruction (RJR) technique. RJR can be considered the conceptual successor
107 of the razor variables. We impose a particular final state “decay tree” on an events,
108 which roughly corresponds to a simplified Feynmann diagram in decays containing
109 weakly-interacting particles. We account for the missing degrees of freedom associated

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

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

The Standard Model

2.1 Overview

128 A Standard Model is another name for a theory of the internal symmetry group
 129 $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$, with its associated set of parameters. *The* Standard
 130 Model refers specifically to a Standard Model with the proper parameters to describe
 131 the universe. The SM is the culmination of years of work in both theoretical
 132 and experimental particle physics. In this thesis, we take the view that theorists
 133 construct a model with the field content and symmetries as inputs, and write down the
 134 most general Lagrangian consistent with those symmetries. Assuming this model is
 135 compatible with nature (in particular, the predictions of the model are consistent with
 136 previous experiments), experimentalists are responsible measuring the parameters of
 137 this model This will be applicable for this chapter and the following one.

cite

138 Additional theoretical background is in 9.6. The philosophy and notations are
 139 inspired by [46, 47].

2.2 Field Content

The Standard Model field content is

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

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

144 We observed that Q_L, U_R , and D_R are triplets under $SU(3)_C$; these are the *quark*
 145 fields. The *color* group, $SU(3)_C$ is mediated by the *gluon* field $G^\mu(8, 1)_0$, which has
 146 8 degrees of freedom. The fermion fields $L_L(1, 2)_{-1}$ and $E_R(1, 1)_{-2}$ are singlets under
 147 $SU(3)_C$; we call them the *lepton* fields.

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

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

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

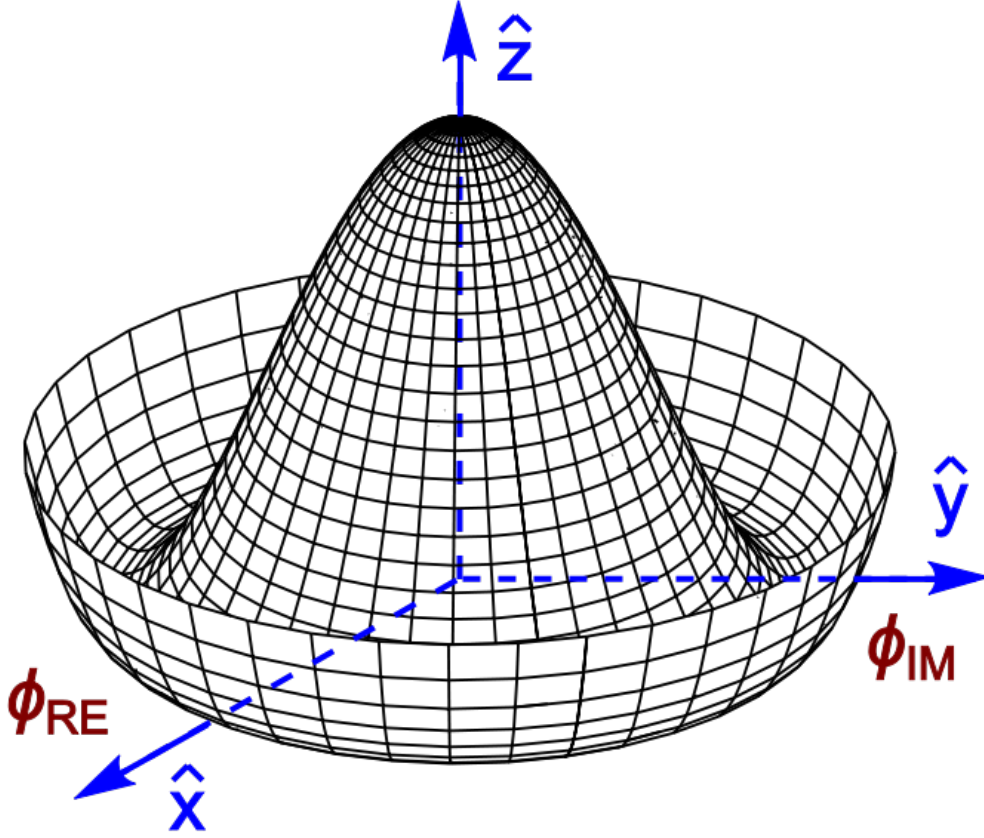
159 **Electroweak sector**

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 ϕ is the complex

Figure 2.1: Sombrero potential



Higgs multiplet. The covariant derivative D^μ 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

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

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

161 The terms in the Lagrangian 2.2 proportional to μ^2 and λ make up the “Higgs
 162 potential” [48]. As normal (see Appendix 9.6), we restrict $\lambda > 0$ to guarantee our
 163 potential is bounded from below, and we also require $\mu^2 < 0$, which gives us the
 164 standard “sombbrero” 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 ϕ to point in the real direction, and write the Higgs field ϕ 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 θ_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

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

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

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

165 We thus see how the Higgs mechanism gives rise to the masses of the W^\pm and Z
 166 boson in the Standard Model; the mass of the photon is zero, as expected. The
 167 $SU(2)_L \otimes U(1)_Y$ symmetry of the initially massless $W_{1,2,3}$ and B fields is broken to
 168 the $U(1)_{EM}$. Of the four degrees of freedom in the complex Higgs doublet, three are
 169 “eaten” when we give mass to the W^\pm and Z_0 , while the other degree of freedom is
 170 the Higgs particle, as found in 2012 by the ATLAS and CMS collaborations [49, 50].

171 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

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

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

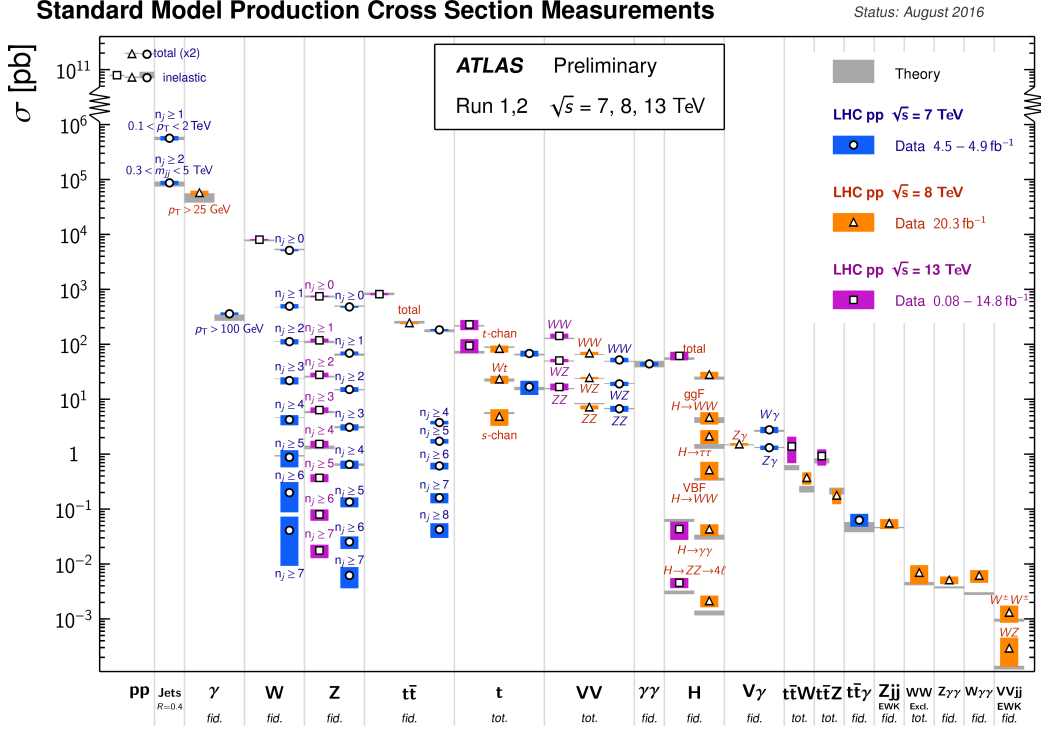
$$G_a^{\mu\nu} = \partial^\mu G_a^\nu - \partial^\nu G_a^\mu - g_s f^{abc} G_b^\mu G_c^\nu, a, b, c = 1, \dots, 8 \quad (2.13)$$

172 where f^{abc} are the structure constants of $SU(3)_C$, which are analogous to ϵ_{abc} for
 173 $SU(2)_L$. The kinetic term for the quarks is contained in the standard ∂_μ term, while
 174 the field strength term contains the interactions between the quarks and gluons, as
 175 well as the gluon self-interactions.

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

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

Figure 2.2: Cross-sections of various Standard Model processes



gluons that interact are part of the *sea* particles inside the proton; the simple $p = uud$ model does not apply. The main *valence* uud quarks are constantly interacting via gluons, which can themselves radiate gluons or split into quarks, and so on. A more useful understanding is given by the colloquially-known *bag* model [51, 52], where the proton is seen as a “bag” of (in principle) infinitely many partons, each with energy $E < \sqrt{s} = 6.5$ TeV. One then collides this (proton) bag with another, and views the products of this very complicated collision, where calculations include many loops in nonperturbative QCD calculations.

Fortunately, we are generally saved by the QCD factorization theorems [53]. This allows one to understand the hard (i.e. short distance or high energy) $2 \rightarrow 2$ parton process using the tools of perturbative QCD, while making series of approximations known as a *parton shower* model to understand the additional corrections from nonperturbative QCD. We will discuss the reconstruction of jets by experiments in Ch.5.

210 Fermions

211 We will now look more closely at the fermions in the Standard Model [54].

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

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

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

215 There is the electron (e), muon (μ), and tau (τ), each of which has an associated
216 neutrino (ν_e, ν_μ, ν_τ). Each of the so-called charged (“electron-like”) leptons has
217 electromagnetic charge -1 , while the neutrinos all have $q_{EM} = 0$.

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

223 As the neutrinos are electrically neutral, nearly massless, and only interact via the
224 weak force, it is quite difficult to observe them directly. Since LHC experiments rely
225 overwhelmingly on electromagnetic interactions to observe particles, the presence of
226 neutrinos is not observed directly. Neutrinos are instead observed by the conservation
227 of four-momentum in the plane transverse to the proton-proton collisions, known as
228 *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 :

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

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

230 Each up-like quark has charge $q_{up} = 2/3$, while the down-like quarks have $q_{down} =$
231 $-1/3$. At the high energies of the LHC, one often makes the distinction between
232 the light quarks (u, d, c, s), the bottom quark, and top quark. In general, due to
233 the hadronization process described above, the light quarks, with masses $m_q < \sim$
234 1.5GeV are indistinguishable by LHC experiments. Their hadronic decay products
235 generally have long lifetimes and they are reconstructed as jets.¹ The bottom quark
236 hadronizes primarily through the B -mesons, which generally travels a short distance
237 before decaying to other hadrons. This allows one to distinguish decays via b -quarks
238 from other jets; this procedure is known as *b-tagging* and will be discussed more in
239 Ch.5. Due to its large mass, the top quark decays before it can hadronize; there
240 are no bound states associated to the top quark. The top is of particular interest at
241 the LHC; it has a striking signature through its most common decay mode $t \rightarrow Wb$.
242 Decays via tops, especially $t\bar{t}$ are frequently an important signal decay mode, or an
243 important background process.

244 Interactions in the Standard Model

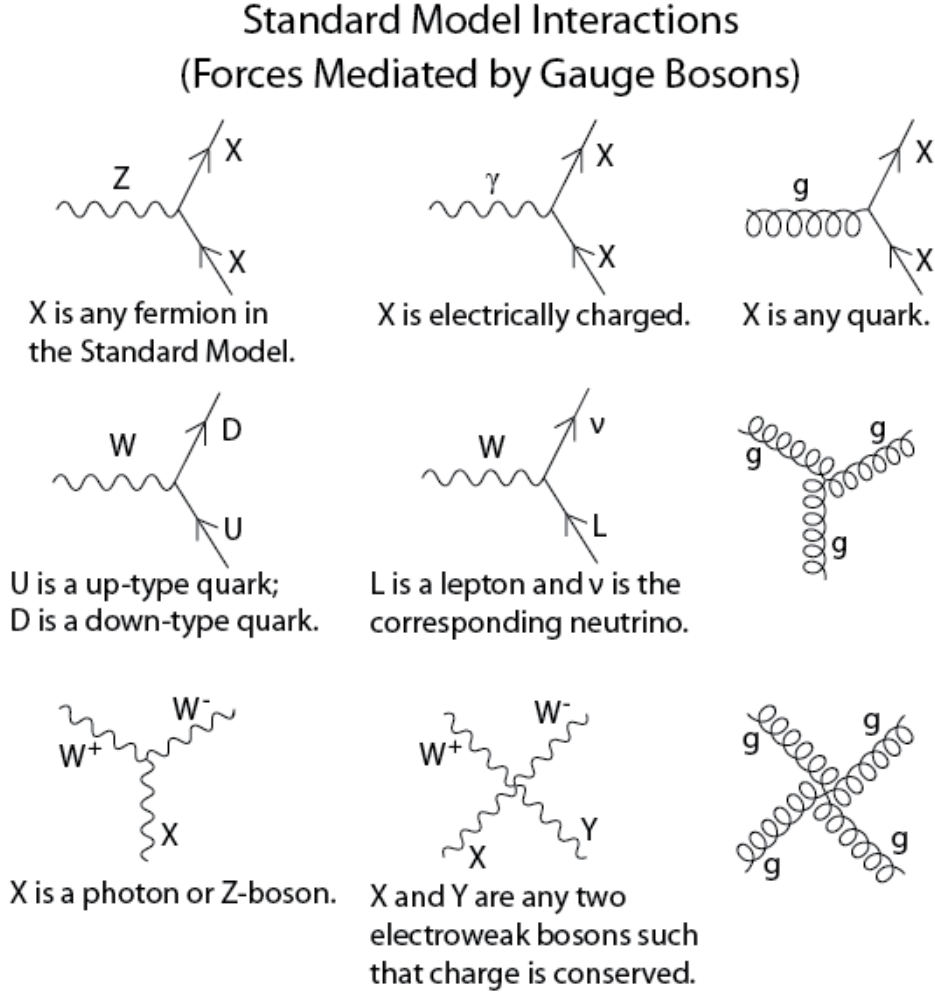
245 We briefly overview the entirety of the fundamental interactions of the Standard
246 Model; these can also be found in 2.3.

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

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

¹In 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



253 mass. The W^\pm has two important three-point interactions with fermions. First, the
 254 W^\pm can interact with an up-like quark and a down-like quark; an important example
 255 in LHC experiments is $t \rightarrow Wb$. The coupling constants for these interactions are
 256 encoded in the unitary matrix known as the Cabibbo–Kobayashi–Maskawa (CKM)
 257 matrix [55, 56], and are generally known as flavor-changing interactions. Secondly,
 258 the W^\pm interacts with a charged lepton and its corresponding neutrino. In this case,
 259 the unitary matrix that corresponds to CKM matrix for quarks is the identity matrix,
 260 which forbids (fundamental) vertices such as $\mu \rightarrow We$. For leptons, instead this is
 261 a two-step process : $\mu \rightarrow \nu_\mu u W \rightarrow \nu_\mu u \bar{\nu}_e e$. Finally, there are the self-interactions

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

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

2.3 Deficiencies of the Standard Model

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

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

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

An additional issue is the lack of *gauge coupling unification*. The couplings of 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 [57–60] and modified minimal subtraction scheme with $m_{\bar{M}S}$ as indicated in the table[61]

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

energy scales) of the theory. The idea is closely related to the unification of the electromagnetic and weak forces at the so-called *electroweak scale* of $O(100 \text{ GeV})$. One would hope this behavior was repeated between the electroweak forces and the strong force at some suitable energy scale. The Standard Model does automatically 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 dependent on the scale of the ultraviolet physics, Λ . 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.

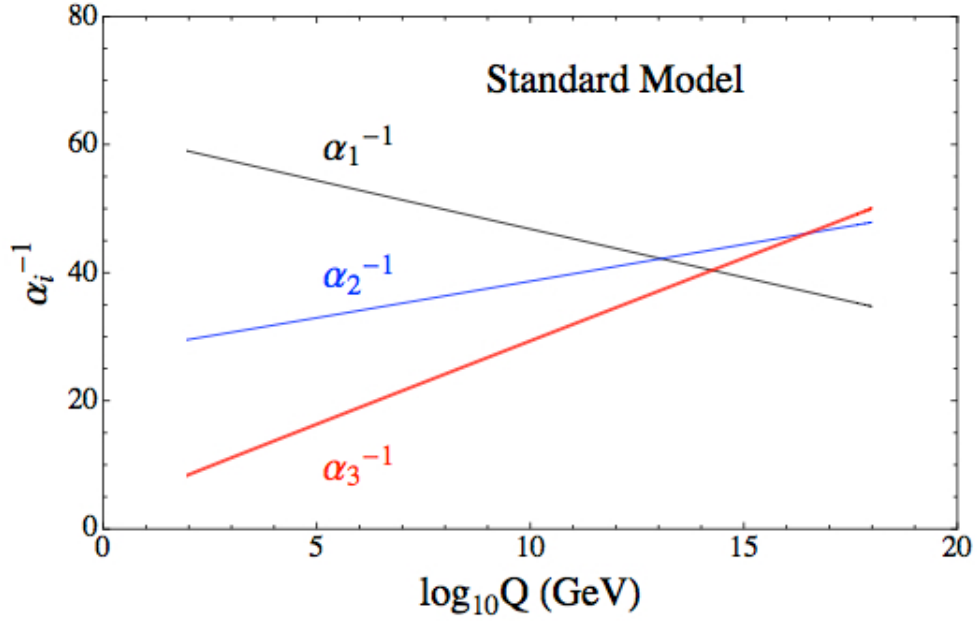
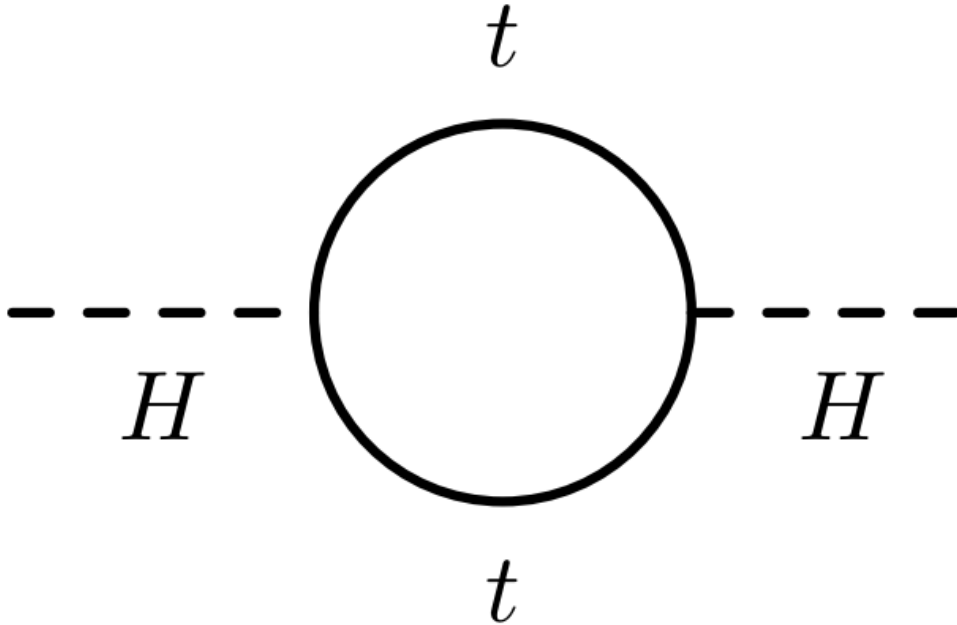


Figure 2.5: The dominant quantum loop correction to the Higgs mass in the Standard Model.



case, we expect the corrections to the Higgs mass like

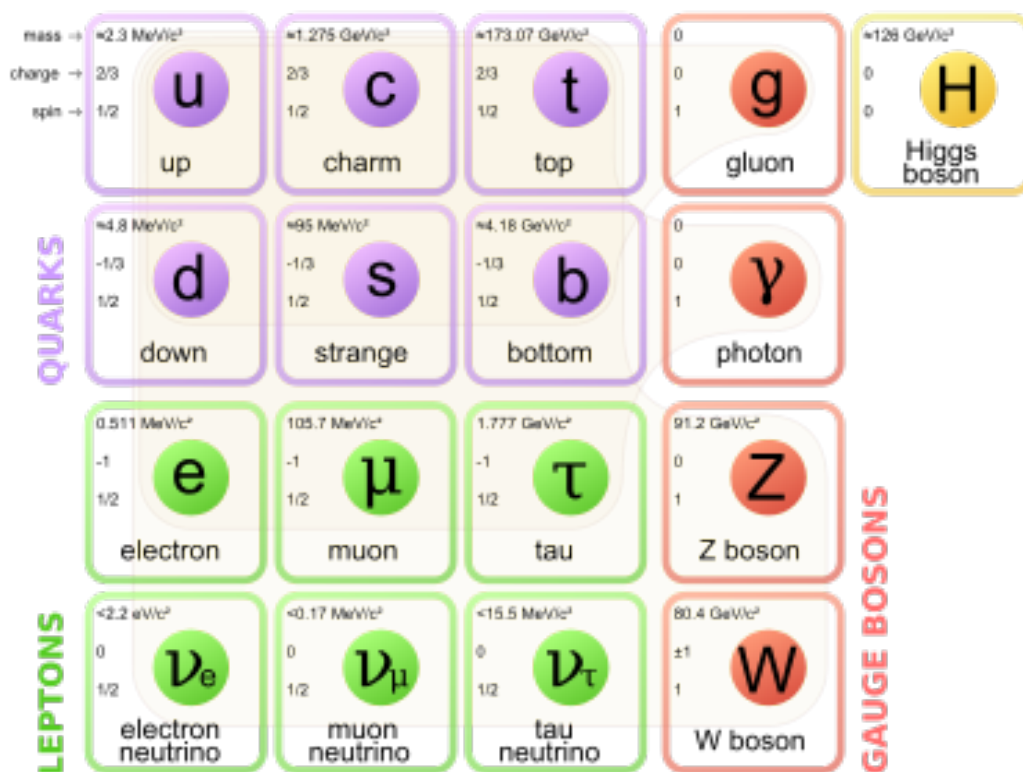
$$\delta m_H^2 \approx \left(\frac{m_t}{8\pi^2 \langle \phi \rangle_{VEV}} \right)^2 \Lambda_{Planck}^2. \quad (2.17)$$

287 To achieve the miraculous cancellation required to get the observed Higgs mass of
 288 125 GeV, one needs to then set the bare Higgs mass m_0 , our input to the Standard
 289 Model Lagrangian, itself to a *precise* value $\sim 10^{19}$ GeV. This extraordinary level of
 290 parameter finetuning is quite undesirable, and within the framework of the Standard
 291 Model, there is little that can be done to alleviate this issue.

292 An additional concern, of a different nature, is the lack of a *dark matter* candidate
 293 in the Standard Model. Dark matter was discovered by observing galactic rotation
 294 curves, which showed that much of the matter that interacted gravitationally was
 295 invisible to our (electromagnetic) telescopes [16–22]. The postulation of the existence
 296 of dark matter, which interacts at least through gravity, allows one to understand
 297 these galactic rotation curves. Unfortunately, no particle in the Standard Model could
 298 possibly be the dark matter particle. The only candidate truly worth another look is
 299 the neutrino, but it has been shown that the neutrino content of the universe is simply
 300 too small to explain the galactic rotation curves [22, 62]. The experimental evidence
 301 from the galactic rotations curves thus show there *must* be additional physics beyond
 302 the Standard Model, which is yet to be understood.

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

Figure 2.6: Particles of the Standard Model



305

Chapter 3

306

Supersymmetry

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

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

310 **3.1 Motivation**

311 **Only Additional allowed Lorentz invariant symmetry**

312 **Dark Matter**

313 **Cancellation of quadratic divergences in corrections to the**

314 **Higgs Mass**

315 **3.2 Supersymmetry**

316 **3.3 Additional particle content**

317 **3.4 Phenomenology**

318 **R parity Consequences for sq/gl decays**

319

Chapter 4

320

The Large Hadron Collider

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

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

324 **4.1 Magnets**

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

328

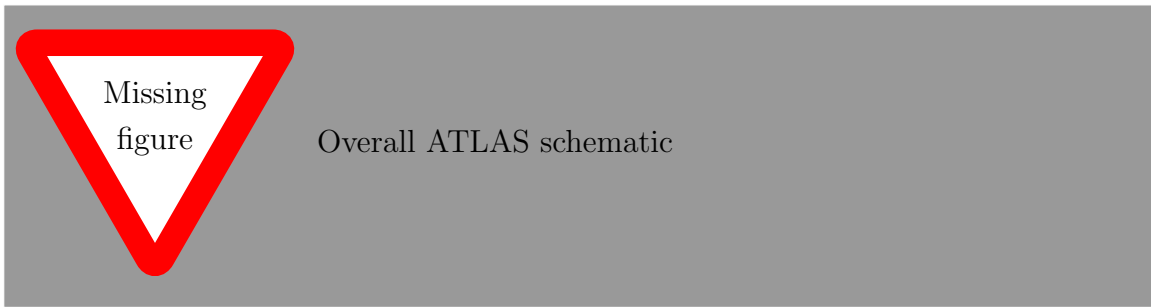
Chapter 5

329

The ATLAS detector

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

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



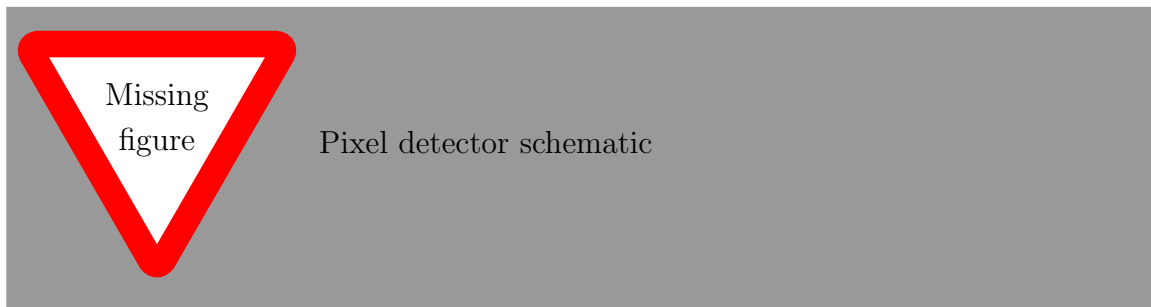
333

334

335 **5.1 Inner Detector**

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337 table of contents. If you want your sections to be numbered and to appear in the
338 table of contents, remove the asterisk.

339 **Pixel Detector**

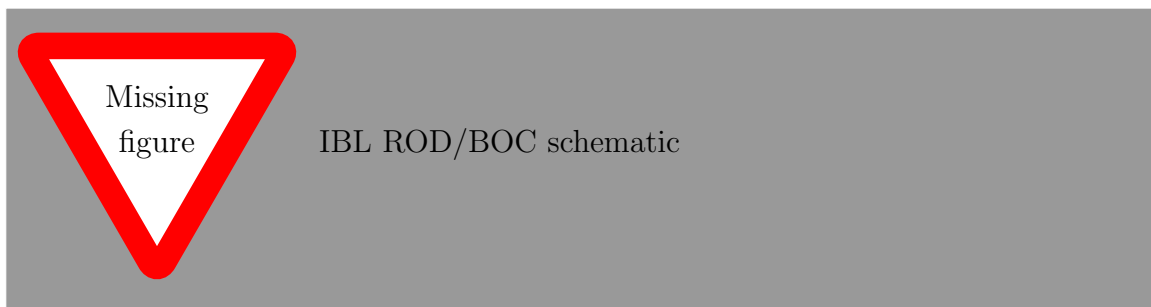


340

341

342 **Insertable B-Layer**

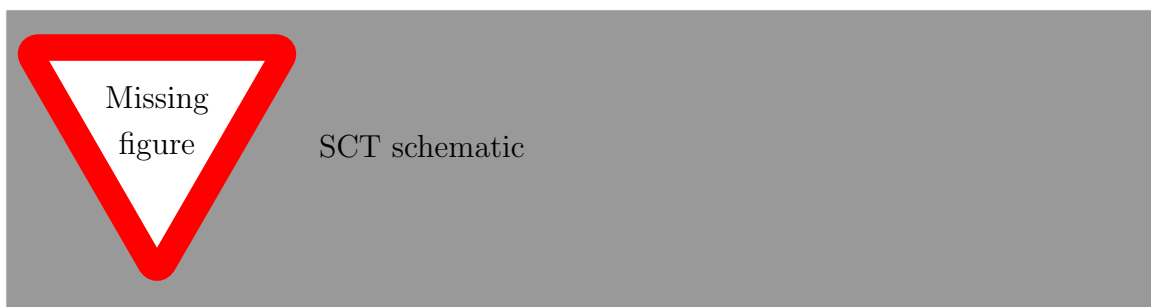
343 Qualification task, so add a bit more.



344

345

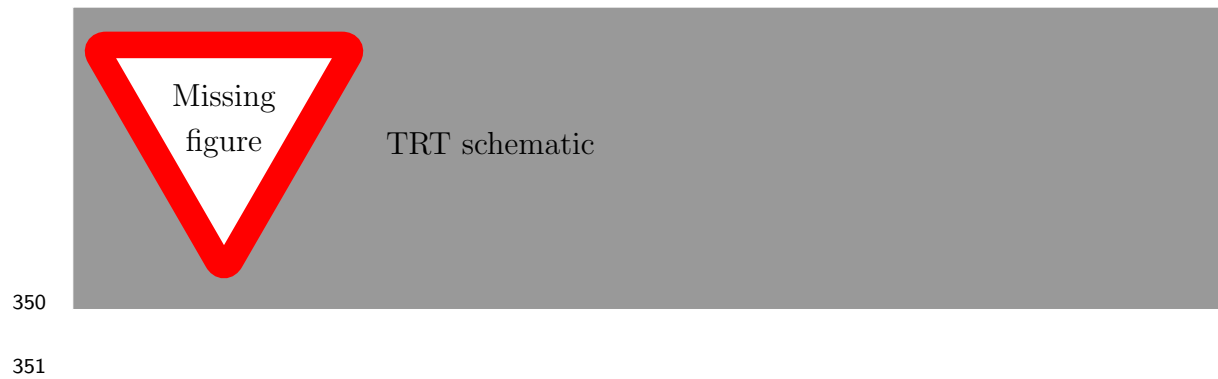
346 **Semiconductor Tracker**



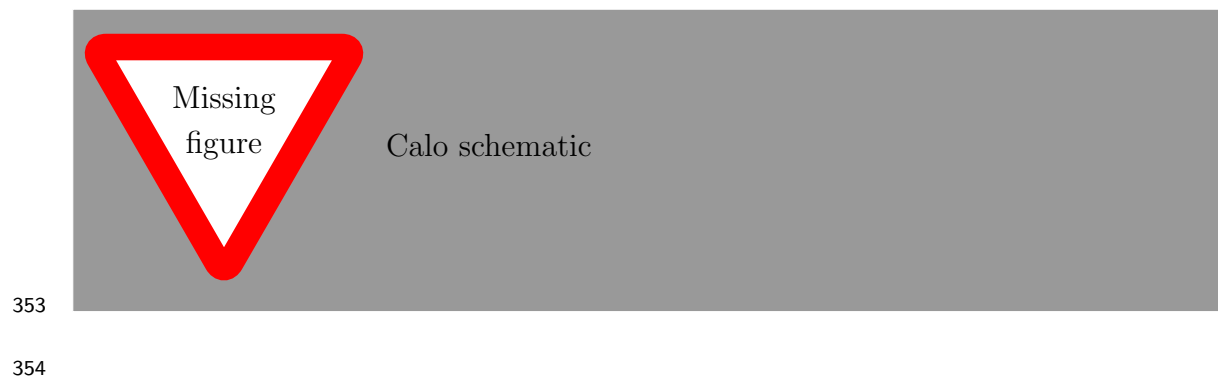
347

348

349 **Transition Radiation Tracker**



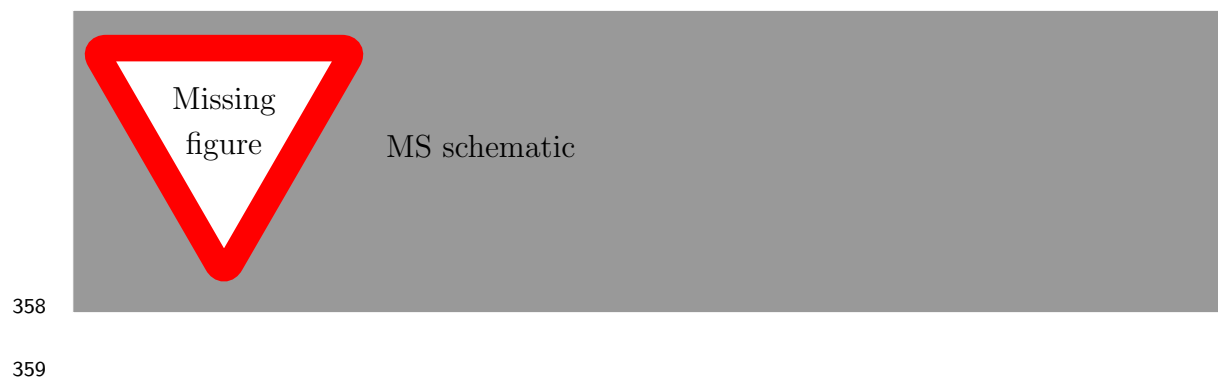
352 **5.2 Calorimeter**



355 **Electromagnetic Calorimeter**

356 **Hadronic Calorimeter**

357 **5.3 Muon Spectrometer**



360

Chapter 6

361

The Recursive Jigsaw Technique

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

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

365 **6.1 Razor variables**

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

369 **6.2 SuperRazor variables**

370 **6.3 The Recursive Jigsaw Technique**

371 **6.4 Variables used in the search for zero lepton**

372 **SUSY**

Title of Chapter 1

375

Chapter 8

376

Title of Chapter 1

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

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

380 **8.1 Object reconstruction**

381 **Photons, Muons, and Electrons**

382 **Jets**

383 **Missing transverse momentum**

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

385 **8.2 Signal regions**

386 **Gluino signal regions**

387 **Squark signal regions**

388 **Compressed signal regions**

389 **8.3 Background estimation**

390 **Z $\nu\nu$**

391 **W $e\nu$**

392 **$t\bar{t}$ bar**

393

Chapter 9

394

Title of Chapter 1

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

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

398 **9.1 Statistical Analysis**

399 maybe to be moved to an appendix

400 **9.2 Signal Region distributions**

401 **9.3 Pull Plots**

402 **9.4 Systematic Uncertainties**

403 **9.5 Exclusion plots**

404

Conclusion

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

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

408 **9.6 New Section**

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

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The Standard Model

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

Quantum Field Theory

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

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

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

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

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

cite Yuval's
lectures
and notes
somehow

cite

- 594 3. Reality condition - The Lagrangian is real to conserve probability.
- 595 4. Lorentz invariance - The Lagrangian is invariant under the Poincaré group of
596 spacetime.
- 597 5. Analyticity - The Lagrangian is an analytical function of the fields; this is to
598 allow the use of perturbation theory.
- 599 6. Invariance and Naturalness - The Lagrangian is invariant under some internal
600 symmetry groups; in fact, the Lagrangian will have *all* terms allowed by the
imposed symmetry groups.
maybe add in ref here
- 602 7. Renormalizability - The Lagrangian will be renormalizable - in practice, this
603 means there will not be terms with more than power 4 in the fields.

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

607 Symmetries

608 Symmetries can be seen as the fundamental guiding concept of modern physics.
cite? 609 Symmetries are described by “groups”. . To illustrate the importance of symmetries
610 and their mathematical description, groups, we start here with two of the simplest
611 and most useful examples : \mathbb{Z}_2 and $U(1)$.

612 \mathbb{Z}_2 symmetry

613 \mathbb{Z}_2 symmetry is the simplest example of a “discrete” symmetry. Consider the most
614 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

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

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

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

619 The effect of this symmetry is that the total number of ϕ particles can only change
 620 by even numbers, since the only interaction term $\lambda\phi^4$ is an even power of the field.
 621 This symmetry is often imposed in supersymmetric theories, as we will see in Chapter
 622 3.

623 **$U(1)$ symmetry**

624 $U(1)$ is the simplest example of a continuous (or *Lie*) group. Now consider a theory
 625 with a single complex scalar field $\phi = \text{Re } \phi + i \text{Im } \phi$

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

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

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

Local symmetries

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

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

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

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

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

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

$$(9.10)$$

Since ϕ and ϕ^* transform with the opposite phase, this will lead to the invariance of the Lagrangian under our local gauge transformation. This D^μ is of the following form

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

644 where A^μ is a vector field we introduce with the transformation law

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

645 and g is the coupling constant associated to vector field. This vector field A^μ is
 646 also known as a “gauge” field.

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

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

648 and then we must also add the kinetic term :

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

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

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

651 Symmetry breaking and the Higgs mechanism

652 Here we view some examples of symmetry breaking. We investigate breaking of a
 653 global $U(1)$ symmetry and a local $U(1)$ symmetry. The SM will break the electroweak
 654 symmetry $SU(2) \times U(1)$, and in Chapter 3 we will see how supersymmetry must also
 655 be broken.

656 There are two ideas of symmetry breaking

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

660 $\mu \ll \lambda$. In this case, we can often ignore the small term when considering
661 low-energy processes.

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

667 Symmetry breaking a

668 **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 ξ : $\phi = (h + i\xi)/\sqrt{2}$.

The Lagrangian can then be written as

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

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 ϕ , we can see that the “vacuum expectation value” of the theory is

$$2 \langle \phi^\dagger \phi \rangle = \langle h^2 + \xi^2 \rangle = v^2 \quad (9.20)$$

669 We now reach the “breaking” point of this procedure. In the (h, ξ) plane, the
670 minima form a circle of radius v . We are free to choose any of these minima to expand
671 our Lagrangian around; the physics is not affected by this choice. For convenience,
672 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

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