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A search for sparticles in zero lepton final states

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ABSTRACT

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A search for sparticles in zero lepton final states

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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.

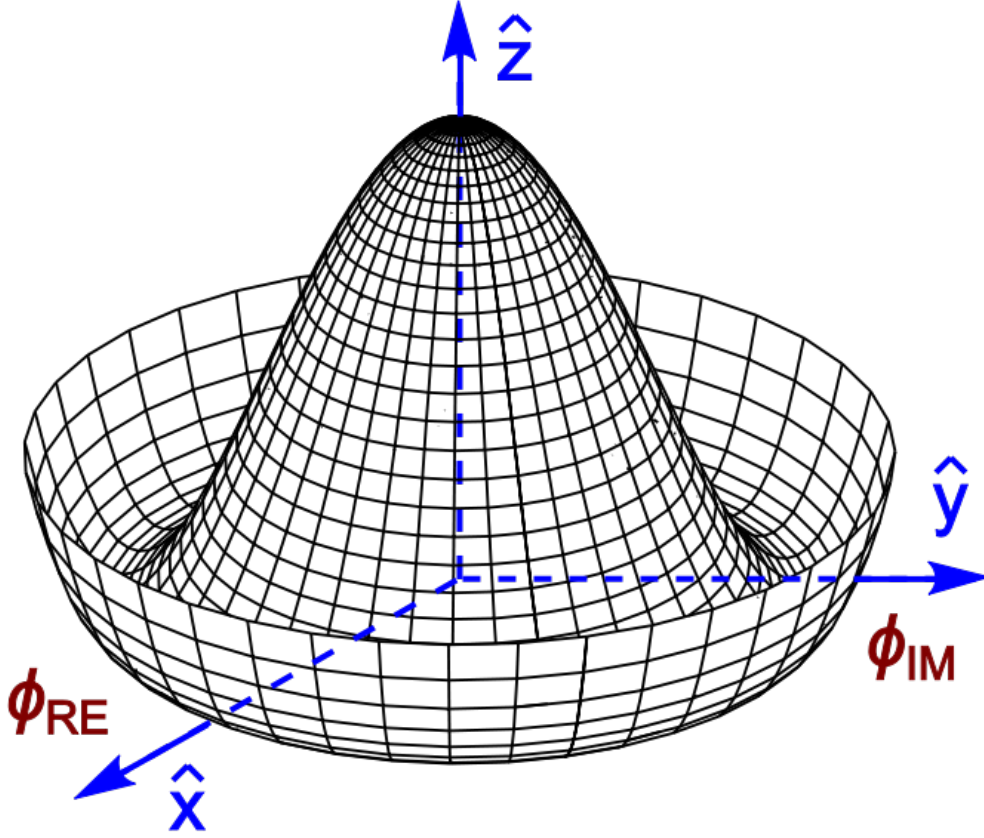
159 **Electroweak sector**

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

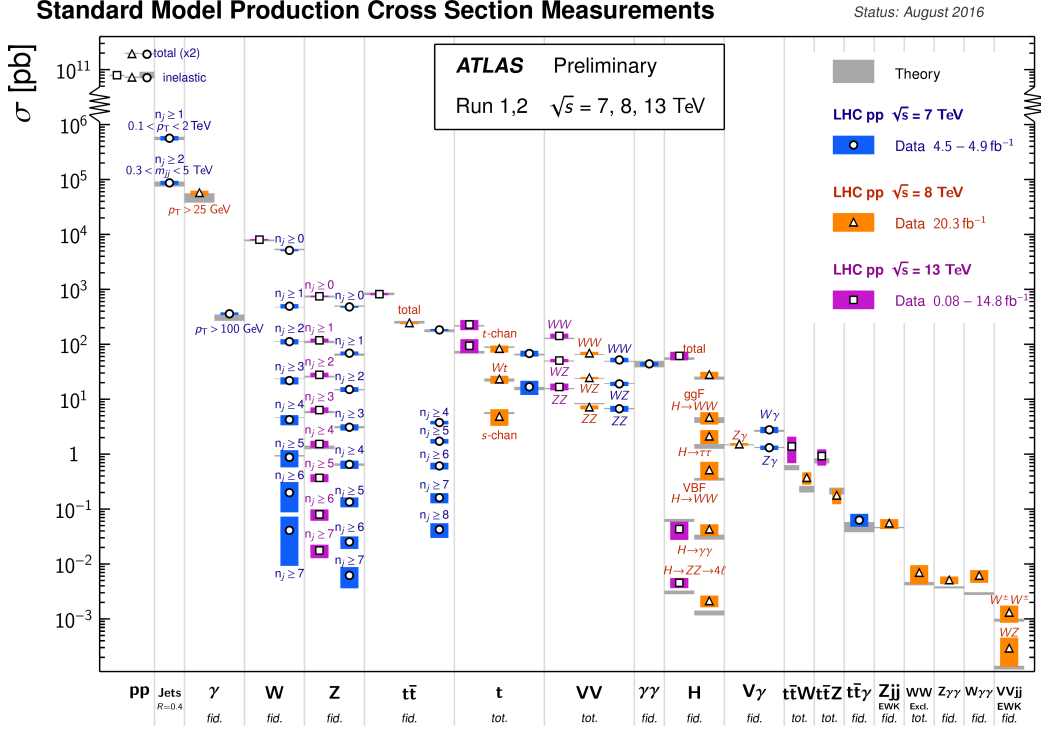
$$G_a^{\mu\nu} = \partial^\mu G_a^\nu - \partial^\nu G_a^\mu - g_s f^{abc} G_b^\mu G_c^\nu, a, b, c = 1, \dots, 8 \quad (2.13)$$

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



196 gluons that interact are part of the *sea* particles inside the proton; the simple $p = uud$
 197 model does not apply. The main *valence* uud quarks are constantly interacting via
 198 gluons, which can themselves radiate gluons or split into quarks, and so on. A more
 199 useful understanding is given by the colloquially-known *bag* model [51, 52], where the
 200 proton is seen as a “bag” of (in principle) infinitely many partons, each with energy
 201 $E < \sqrt{s} = 6.5$ TeV. One then collides this (proton) bag with another, and views the
 202 products of this very complicated collision, where calculations include many loops in
 203 nonperturbative QCD calculations.

204 Fortunately, we are generally saved by the QCD factorization theorems [53]. This
 205 allows one to understand the hard (i.e. short distance or high energy) $2 \rightarrow 2$ parton
 206 process using the tools of perturbative QCD, while making series of approximations
 207 known as a *parton shower* model to understand the additional corrections from
 208 nonperturbative QCD. We will discuss the reconstruction of jets by experiments in
 209 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

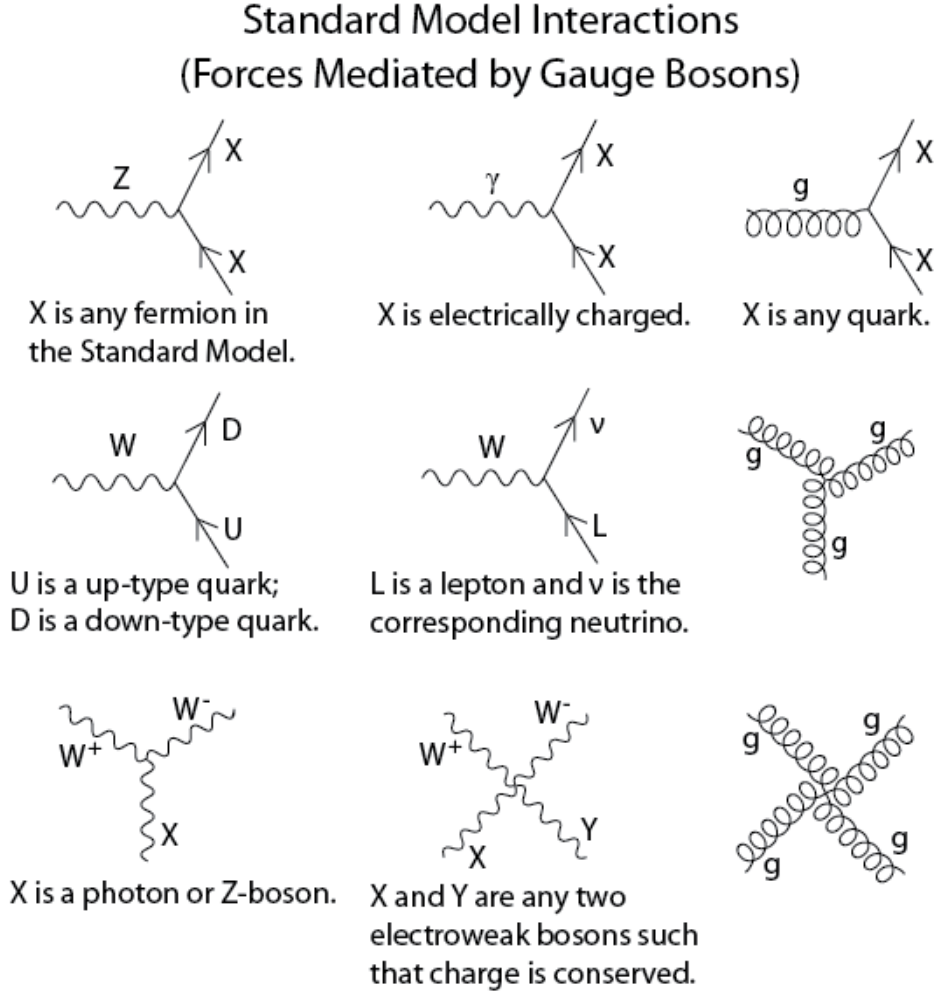
245 We briefly overview the entirety of the fundamental interactions of the Standard
246 Model; these can also be found in 2.3.

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.

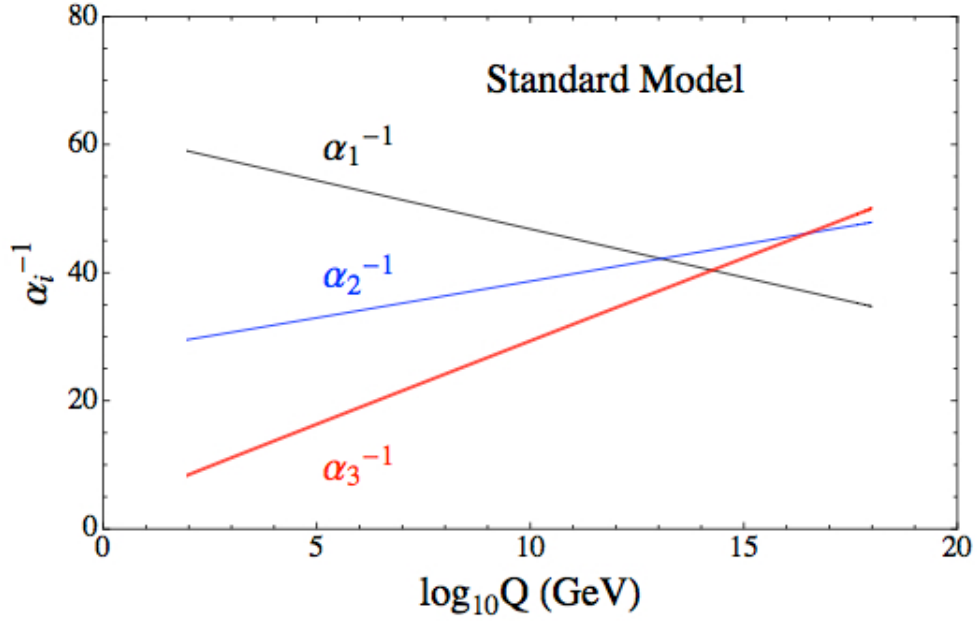
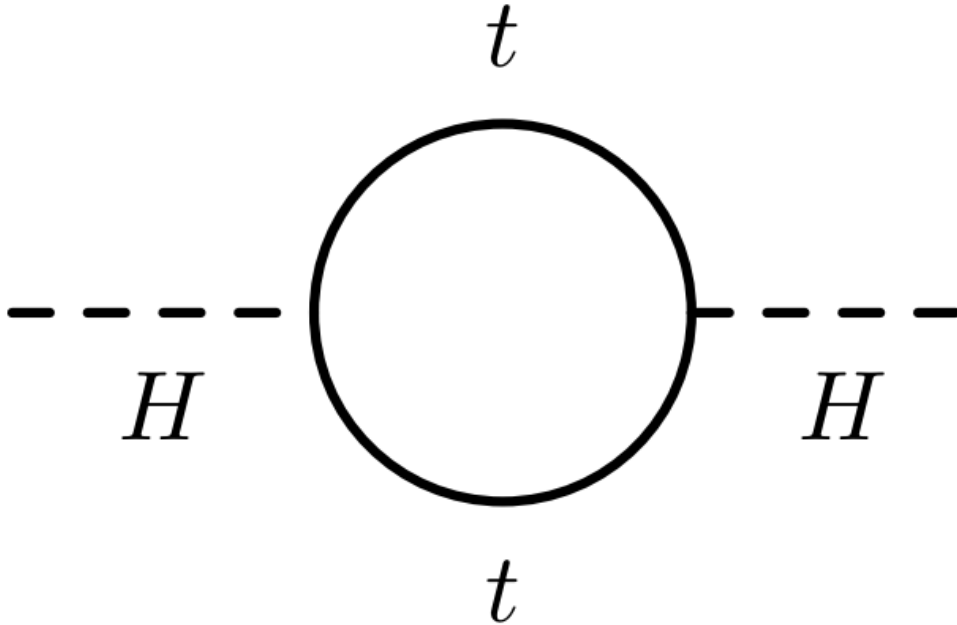


Figure 2.5: The dominant quantum loop correction to the Higgs mass in the Standard Model.



case, we expect the corrections to the Higgs mass like

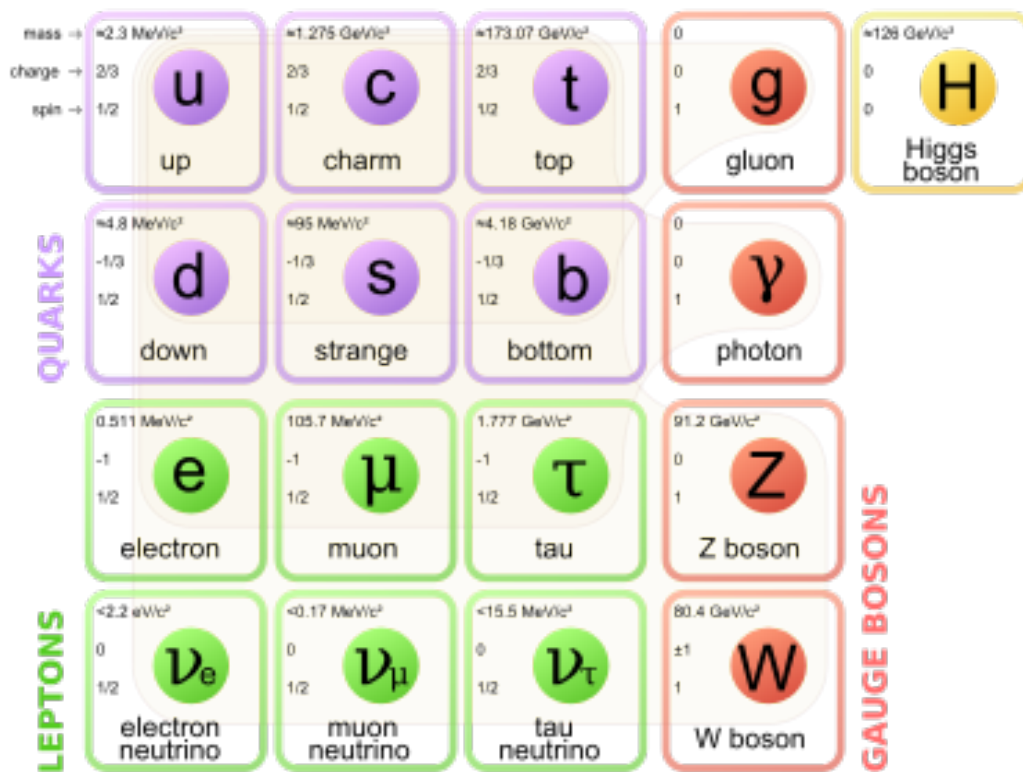
$$\delta m_H^2 \approx \left(\frac{m_t}{8\pi^2 \langle \phi \rangle_{VEV}} \right)^2 \Lambda_{Planck}^2. \quad (2.17)$$

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 This chapter will introduce supersymmetry (SUSY) [15, 63] .. We will begin by
 308 introducing the concept of a *superspace*, and discuss some general ingredients of
 309 supersymmetric theories. This will include a discussion of how the problems with the
 310 Standard Model described in Ch.2 are naturally fixed by these theories.

cite more

cite SUSY

lectures from

presusy

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

3.1 Supersymmetric theories : from space to superspace

Coleman-Mandula “no-go” theorem

319 We begin the theoretical motivation for supersymmetry by citing the “no-go” theorem
 320 of Coleman and Mandula [64]. This theorem forbids *spin-charge unification*; it
 321 states that all quantum field theories which contain nontrivial interactions must be
 322 a direct product of the Poincaré group of Lorentz symmetries, the internal product
 323 from of gauge symmetries, and the discrete symmetries of parity, charge conjugation,
 324 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, 65]. 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* [66]. We will not investigate this view in detail, but it is also a quite intuitive and beautiful way to construct supersymmetry[15].

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 [65] of the Coleman-Mandula theorem [64] is quite restrictive about the forms of such a symmetry. Here, we simply write the (anti-) commutation relations [15] :

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

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

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

334 Supermultiplets

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

347 Since each supermultiplet must contain a fermion state, the simplest type of
348 supermultiplet contains a single Weyl fermion state ($n_F = 2$) which is paired with
349 $n_B = 2$ scalar bosonic degrees of freedom. This is most conveniently constructed as
350 single complex scalar field. We call this construction a *scalar supermultiplet* or *chiral*
351 *supermultiplet*. The second name is indicative; only chiral supermultiplets can contain
352 fermions whose right-handed and left-handed components transform differently under
353 the gauge interactions (as of course happens in the Standard Model).

354 The second type of supermultiplet we construct is known as a *gauge* supermul-
355 tiplet. We take a spin-1 gauge boson (which must be massless due to the gauge
356 symmetry, so $n_B = 2$) and pair this with a single massless Weyl spinor¹. The gauge
357 bosons transform as the adjoint representation of their respective gauge groups;
358 their fermionic partners, which are known as gauginos, must also. In particular,
359 the left-handed and right-handed components of the gaugino fermions have the same

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

360 gauge transformation properties.

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

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

370 3.2 Minimally Supersymmetric Standard Model

cite 371 To construct what is known as the MSSM, we need a few ingredients and assumptions.
372 First, we match the Standard Model particles with their corresponding superpartners
373 of the MSSM. We will also introduce the naming of the superpartners (also known as
374 *sparticles*). We discuss a very common additional restraint imposed on the MSSM,
375 known as R -parity. We also discuss the concept of soft supersymmetry breaking and
376 how it manifests itself in the MSSM.

377 Chiral supermultiplets

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

384 (for “scalar quarks”, “scalar leptons”, and “scalar fermion”²). The “s-” prefix
385 can also be added to the individual quarks i.e. *selectron*, *sneutrino*, and *stop*. The
386 notation is to add a \sim over the corresponding Standard Model particle i.e. \tilde{e} , the
387 selectron is the superpartner of the electron. The two-component Weyl spinors of the
388 Standard Model must each have their own (complex scalar) partner i.e. e_L, e_R have
389 two distinct partners : \tilde{e}_L, \tilde{e}_R . As noted above, the gauge interactions of any of the
390 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:

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

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

$$(3.8)$$

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

²The last one should probably have bigger scare quotes.

397 Gauge supermultiplets

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

401 The first gauge supermultiplet contains the gluon, and its superpartner, which is
 402 known as the *gluino*, denoted \tilde{g} . The gluon is of course the SM mediator of $SU(3)_C$;
 403 the gluino is also a colored particle, subject to $SU(3)_C$. From the SM before EWSB,
 404 we have the four gauge bosons of the electroweak symmetry group $SU(2)_L \otimes U(1)_Y$:
 405 $W^{1,2,3}$ and B^0 . The superpartners of these particles are thus the *winos* $\tilde{W}^{1,2,3}$ and
 406 *bino* \tilde{B}^0 , where each is placed in another gauge supermultiplet with its corresponding
 407 SM particle. After EWSB, without breaking supersymmetry, we would also have the
 408 zino \tilde{Z}^0 and photino $\tilde{\gamma}$.

409 At this point, it's important to take a step back. Where are these particles?
 410 As stated above, supersymmetric theories require that the masses and all quantum
 411 numbers of the SM particle and its corresponding sparticle are the same. Of course,
 412 we have not observed a selectron, squark, or wino. The answer, as it often is, is that
 413 supersymmetry is *broken* by the vacuum state of nature [15].

414 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

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

415 where B, L is the baryon (lepton) number and s is the spin. The imposition of
 416 this symmetry forbids certain terms from the MSSM Lagrangian that would violate
 417 baryon and/or lepton number . This is required³ in order to prevent proton decay, as

Feynmann
 diagram

418 shown in Fig..

figure

419 In supersymmetric models, this is a \mathbb{Z}_2 symmetry, where SM particles have $R = 1$
420 and sparticles have $R = -1$. We will take R -parity as part of the definition of
421 the MSSM. We will discuss later the *drastic* consequences of this symmetry on SUSY
422 phenomenology⁴..

do I want
to be this
strong

423 Soft supersymmetry breaking

The fundamental idea of *soft* supersymmetry breaking[15, 63, 67, 68] is that we would like to break supersymmetry without reintroducing the quadratic divergences we discussed at the end of Chapter 2. Assuming we can do this procedure, we can write the Lagrangian in a form :

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

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lectures here
also

something
smarter here

424 In this sense, the symmetry breaking is “soft”, since we have separated out the
425 completely symmetric terms from those soft terms which will not allow the quadratic
426 divergences in .

section

427 The explicitly allowed terms in the soft-breaking Lagrangian are

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lectures

- 428 • Mass terms for the scalar components of the chiral supermultiplets
- 429 • Mass terms for the Weyl spinor components of the gauge supermultiplets
- 430 • Trilinear couplings of scalar components of chiral supermultiplets

³This is the usual story, but it’s actually a bit more complicated. The author has become quite skeptical of this claim.

⁴The author has actually come to the view that people “like” R -parity conservation precisely because it leads to an interesting phenomenology.

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

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

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

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

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

where we have introduced the following notations :

1. M_3, M_2, M_1 are the gluino, wino, and bino masses.
2. a_u, a_d, a_e are complex 3×3 matrices in family space.
3. $m_Q^2, m_u^2, m_d^2, m_L^2, m_e^2$ are hermitian 3×3 matrices in family space.
4. $m_{H_u}^2, m_{H_d}^2, b$ are the SUSY-breaking contributions to the Higgs potential.

We have written matrix terms without any sort of additional notational decoration to indicate their matrix nature, and we now show why. The first term are straightforward; these are just the straightforward mass terms for these fields. There are strong constraints on the off-diagonal terms for the matrices of 2 [69, 70]; for simplicity, we will assume that each $a_i, i = u, d, e$ is proportional to the Yukawa coupling matrix : $a_i = A_{i0} y_i$. The matrices in ?? can be similarly constrained by experiments [70–78] Finally, we assume that the elements 4 contributing to the Higgs potential as well as all of the 1 terms must be real, which limits the possible CP-violating interactions to those of the Standard Model. We thus only consider

flavor-blind, CP-conserving interactions within the MSSM.

3.3 Phenomenology

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

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

add feyn-
mann
diagram
example

For this thesis, we will be presenting an inclusive search for squarks and gluinos with zero leptons in the final state. This is a very interesting decay channel⁵, due to the high cross-sections of $\tilde{g}\tilde{g}$ and $\tilde{q}\tilde{q}$ decays. This is a direct consequence of the fact that these are the colored particles of the MSSM. Since the sparticles interact with the gauge groups of the SM in the same way as their SM partners, the colored sparticles, the squarks and gluinos, are produced and decay as governed by the color group $SU(3)_C$ with the strong coupling g_s . The digluino production is particularly copious, due to color factor corresponding to the color octet of $SU(3)_C$.

Figure

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

For digluino production, the most common decay is $\tilde{g} \rightarrow g\tilde{q}$, due to the large g_s

⁵Prior to Run2, probably the most *most* interesting SUSY decay channel.

coupling. The squark then decays as listed above. In this case, we generically search for four jets and missing transverse energy from the LSPs. We can also have the squark decay in association with a W^\pm ; in this thesis, we are interested in those cases where this W^\pm goes hadronically.

In the context of experimental searches for SUSY, we often consider *simplified models*. These models make certain assumptions which allow easy comparisons of results by theorists and rival experimentalists. In the context of this thesis, the simplified models will make assumptions about the branching ratios described in the preceding paragraphs. In particular, we will often choose a model where the decay of interest occurs with 100% branching ratio. This is entirely for ease of interpretation by other physicists⁶, but it is important to recognize that these are more a useful comparison tool, especially with limits, than a strict statement about the potential masses of sought-after beyond the Standard Model particle.

3.4 How SUSY solves the problems with the SM

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

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

$$\delta m_H^2 \approx \left(\frac{m_t}{8\pi^2 \langle \phi \rangle_{VEV}} \right)^2 \Lambda_{Planck}^2. \quad (3.15)$$

The miraculous thing about SUSY⁷ is each of these terms *automatically* comes with a term which exactly cancels this contribution[15]. The fermions and bosons have

⁶In 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.

opposite signs in this loop diagram to all orders in perturbation theory, which completely solves the hierarchy problem. Although this wasn't the original inspiration for supersymmetry, it fits so neatly, this is the most motivational reason for supersymmetry.

Gauge coupling unification

496

plot

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

reffig:

Dark matter

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

maybe other

some sort
of plot
probably?
WIMP
exclusions

⁷It is often said that SUSY was created to solve the hierarchy problem, but this is really a rewriting of history. SUSY is the only way out of the Coleman-Mandula theorem, and *happens* to solve the hierarchy problem. This is even more beautiful!

512

Chapter 4

513

The Large Hadron Collider

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

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

517 **4.1 Magnets**

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

521

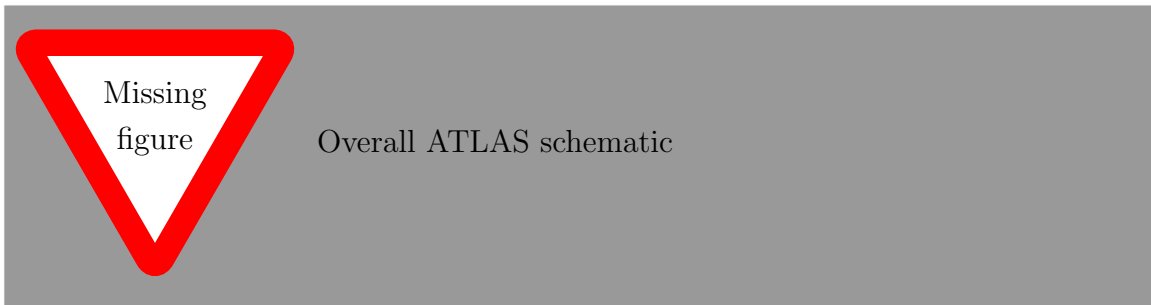
Chapter 5

522

The ATLAS detector

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524 sentence its own line.

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



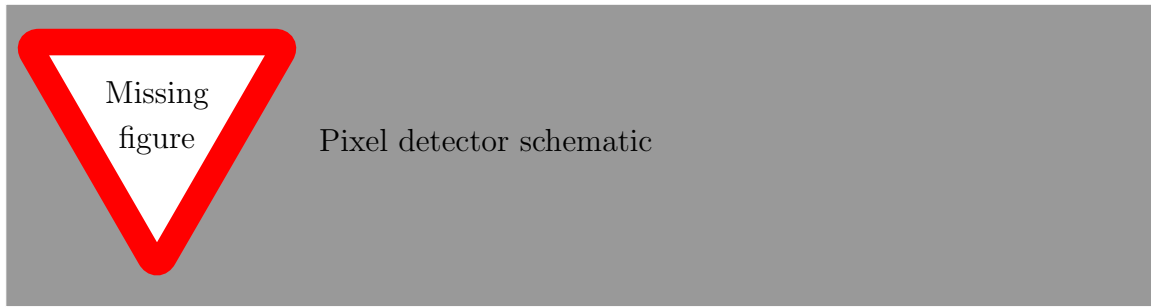
526

527

5.1 Inner Detector

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

532 **Pixel Detector**

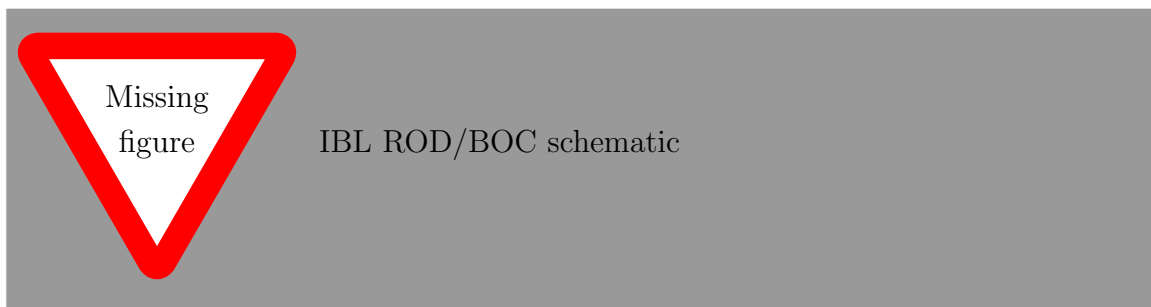


533

534

535 **Insertable B-Layer**

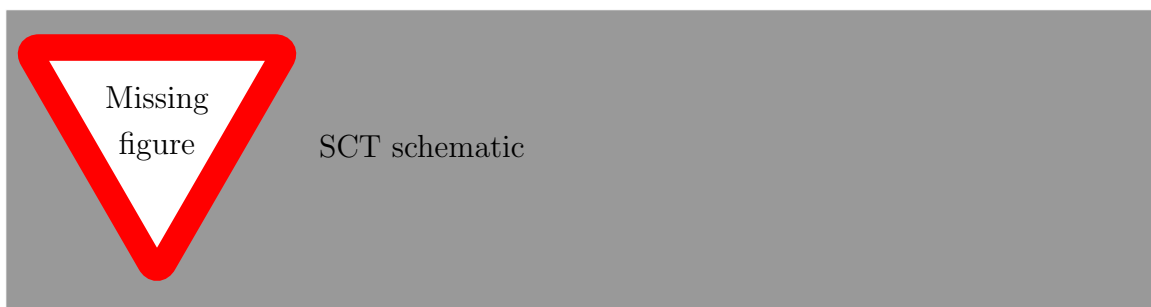
536 Qualification task, so add a bit more.



537

538

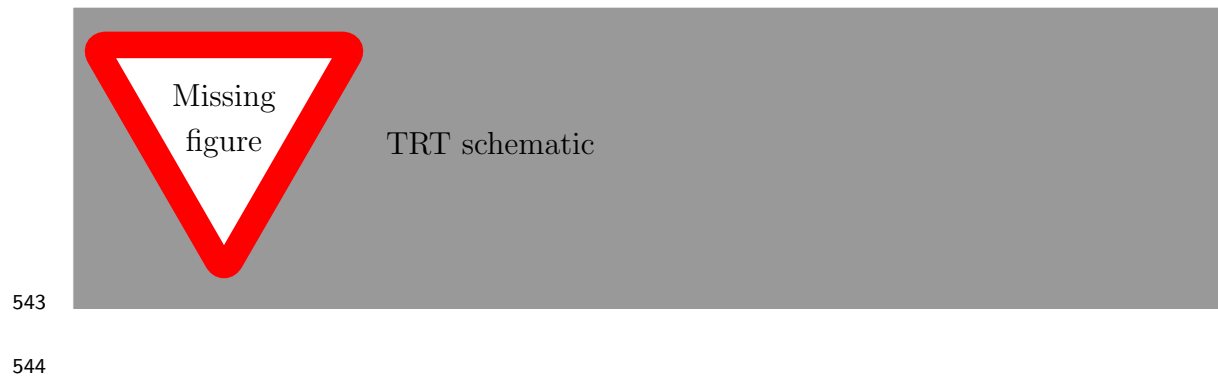
539 **Semiconductor Tracker**



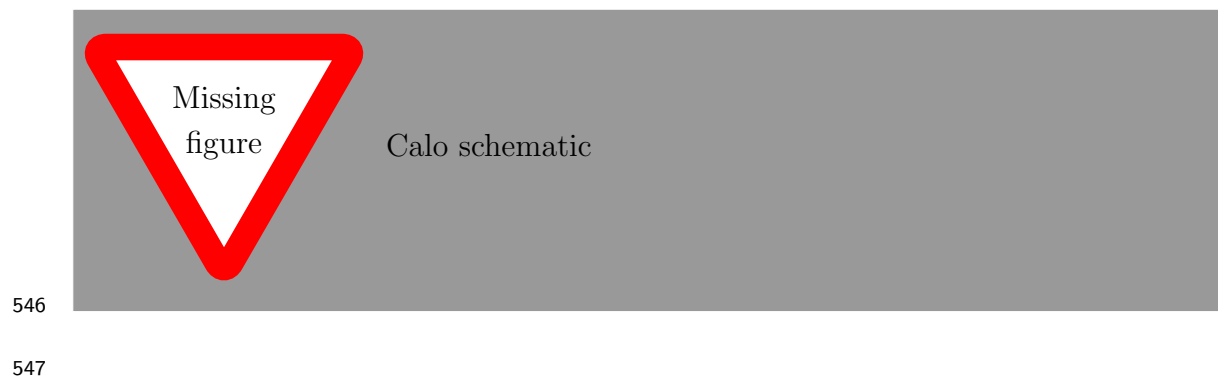
540

541

542 **Transition Radiation Tracker**



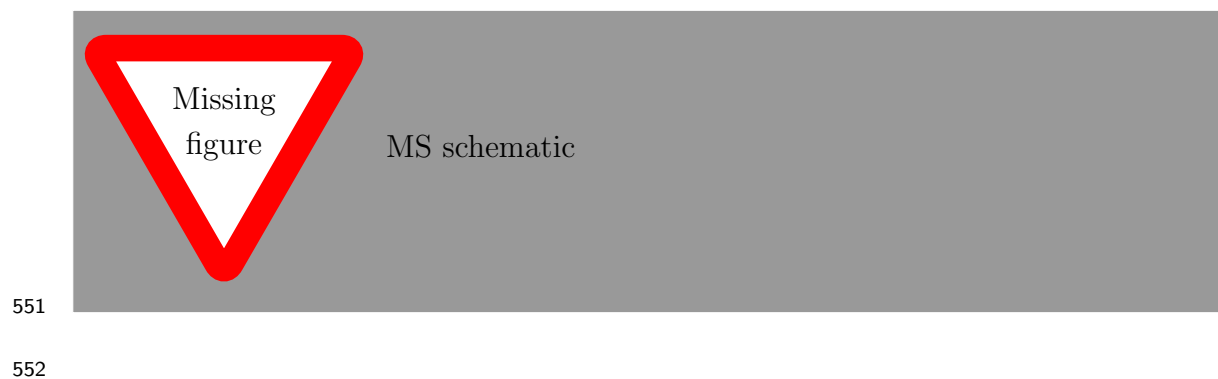
545 **5.2 Calorimeter**



548 **Electromagnetic Calorimeter**

549 **Hadronic Calorimeter**

550 **5.3 Muon Spectrometer**



553

Chapter 6

554

The Recursive Jigsaw Technique

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

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

558 **6.1 Razor variables**

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

562 **6.2 SuperRazor variables**

563 **6.3 The Recursive Jigsaw Technique**

564 **6.4 Variables used in the search for zero lepton**

565 **SUSY**

Title of Chapter 1

568

Chapter 8

569

Title of Chapter 1

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

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

573 **8.1 Object reconstruction**

574 **Photons, Muons, and Electrons**

575 **Jets**

576 **Missing transverse momentum**

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

578 **8.2 Signal regions**

579 **Gluino signal regions**

580 **Squark signal regions**

581 **Compressed signal regions**

582 **8.3 Background estimation**

583 **Z $\nu\nu$**

584 **W $e\nu$**

585 **$t\bar{t}$**

586

Chapter 9

587

Title of Chapter 1

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

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

591 **9.1 Statistical Analysis**

592 maybe to be moved to an appendix

593 **9.2 Signal Region distributions**

594 **9.3 Pull Plots**

595 **9.4 Systematic Uncertainties**

596 **9.5 Exclusion plots**

597

Conclusion

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

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

601 **9.6 New Section**

602 By using the asterisk to start a new section, I keep the section from appearing in the
603 table of contents. If you want your sections to be numbered and to appear in the
604 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.

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- 831 3. Reality condition - The Lagrangian is real to conserve probability.
- 832 4. Lorentz invariance - The Lagrangian is invariant under the Poincaré group of
833 spacetime.
- 834 5. Analyticity - The Lagrangian is an analytical function of the fields; this is to
835 allow the use of perturbation theory.
- 836 6. Invariance and Naturalness - The Lagrangian is invariant under some internal
837 symmetry groups; in fact, the Lagrangian will have *all* terms allowed by the
imposed symmetry groups.
838
839 7. Renormalizability - The Lagrangian will be renormalizable - in practice, this
840 means there will not be terms with more than power 4 in the fields.

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

844 Symmetries

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

849 \mathbb{Z}_2 symmetry

850 \mathbb{Z}_2 symmetry is the simplest example of a “discrete” symmetry. Consider the most
851 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

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

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

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

856 The effect of this symmetry is that the total number of ϕ particles can only change
 857 by even numbers, since the only interaction term $\lambda\phi^4$ is an even power of the field.
 858 This symmetry is often imposed in supersymmetric theories, as we will see in Chapter
 859 3.

860 **$U(1)$ symmetry**

861 $U(1)$ is the simplest example of a continuous (or *Lie*) group. Now consider a theory
 862 with a single complex scalar field $\phi = \text{Re } \phi + i \text{Im } \phi$

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

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

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

866 Local symmetries

867 The two examples considered above are “global” symmetries in the sense that the
868 symmetry transformation does not depend on the spacetime coordinate x_μ . We know
869 look at local symmetries; in this case, for example with a local $U(1)$ symmetry, the
870 transformation has the form $\phi(x_\mu) \rightarrow e^{i\theta(x_\mu)}\phi(x_\mu)$. These symmetries are also known
871 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
874 under a gauge symmetry. This would lead to a model with no dynamics, which is
875 clearly unsatisfactory.

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

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

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

$$(9.10)$$

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

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

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

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

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

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

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

885 and then we must also add the kinetic term :

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

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

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

888 Symmetry breaking and the Higgs mechanism

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

893 There are two ideas of symmetry breaking

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

897 $\mu \ll \lambda$. In this case, we can often ignore the small term when considering
898 low-energy processes.

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

904 Symmetry breaking a

905 **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

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

906 We now reach the “breaking” point of this procedure. In the (h, ξ) plane, the
 907 minima form a circle of radius v . We are free to choose any of these minima to expand
 908 our Lagrangian around; the physics is not affected by this choice. For convenience,
 909 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

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