2

What lies beyond the Standard Model?

The Standard Model is a consistent, renormalizable quantum field theory that accounts for a wide variety of experimental data over an energy range that encompasses a fraction of an electron volt to about 100 GeV, a range of over twelve orders of magnitude.[1] Initially, the SM was tested at the tree level, but the remarkable agreement between SM predictions and the precision measurements at the CERN LEP collider have tested the SM to at least a part per mille and, more importantly, have established that radiative corrections as given by the SM are essential for agreement with these data. Quite aside from this, the SM also qualitatively explains why baryon and lepton numbers appear to be approximately conserved: with the particle content of the SM, it is not possible to write renormalizable interactions that do not conserve baryon and lepton numbers, so that these interactions (if they exist) must be suppressed by (powers of) some new physics scale.

The SM is nevertheless incomplete. Experimental arguments in support of this are:

- **E1** The solar and atmospheric neutrino data, interpreted as neutrino oscillations, strongly suggest neutrinos have mass.
- **E2** Observations, starting with Zwicky in 1933 and continuing to this day with studies of the fluctuations in the spectrum of the relic microwave background from the Big Bang, have established the existence of cold dark matter in the Universe for which there is no candidate in the SM.[2]
- **E3** Observations of type Ia supernovae at large red shifts as well as the cosmic microwave background radiation both suggest that the bulk of the energy of

[1] This is a rather conservative estimate. For instance, it may be argued from the dipole nature of planetary magnetic fields that (at least) electrodynamics has been tested out to Solar System distance scales.

[2] Gravitational microlensing data disfavor black holes (at least in our galactic halo) as dark matter, while dark matter in the form of ordinary baryons condensed into brown dwarfs is excluded, both because it would lead to conflicts with Big Bang nucleosynthesis as well as with the baryon density as determined from the acoustic peaks in the cosmic microwave background spectrum.

the Universe resides in a novel form dubbed “dark energy”. This could be the cosmological constant first introduced by Einstein, or something more bizarre.
- **E4** Gravity exists.

There are also theoretical or aesthetic considerations that suggest that the SM cannot be the complete story.

- **T1** We lack any understanding of particle masses and mixing patterns, which results in the large number of underlying parameters in the SM.
- **T2** The choice of gauge group and particle representations is completely ad hoc.
- **T3** Although we can incorporate the spontaneous breakdown of electroweak symmetry by introducing new scalar fields, we have to do so “by hand” via an arbitrary scalar potential; i.e. there is no understanding of why the squared mass parameter for the Higgs field is negative. Indeed, it remains to be seen whether VEVs of elementary scalar fields are the origin of electroweak symmetry breaking.

While these arguments (especially **E1–E3**) all point to new physics, without further assumptions, they do not point decisively to the scale for this new physics. Fortunately, there is a somewhat different argument that not only suggests that there should be new physics, but also that the scale of the new physics is close to the electroweak scale. To understand this, we must first examine the divergence structure of the radiative corrections in the SM.

2.1 Scalar fields and quadratic divergences

Let us begin by considering radiative corrections to quantum electrodynamics described by the Lagrangian (1.1a). The theory describes the interactions of a fermion with a photon. As discussed in the exercise below, these interactions conserve *chirality*: a left-handed fermion, when it emits (or absorbs) a photon, remains left handed, while a right-handed fermion remains right handed. The kinetic energy term also conserves chirality. Indeed the only term in (1.1a) that does not conserve chirality is the mass term. This observation has an immediate consequence. Because the emission or absorption of a photon cannot change the chirality of the fermion, any radiative correction to the fermion mass (which by the exercise below is an operator that connects ψ_L with ψ_R) must vanish to all orders in perturbation theory if the fermion mass is zero! In other words,

$$\delta m \propto m .$$

It is well known that the loop integrals that enter these calculations are divergent. If we regularize these using a Lorentz invariant cut-off Λ, by dimensional analysis

https://doi.org/10.1017/9781009289801.003 Published online by Cambridge University Press

the cut-off dependence has to be given by,

$$\delta m \propto m \ln \frac{\Lambda}{m}.$$

Naive dimensional analysis would have suggested that $\delta m \propto \Lambda$. However, because of chiral symmetry, the actual divergence is much milder. We say that chiral symmetry protects fermion masses from large radiative corrections. This terminology may seem strange since the correction diverges when we take the cut-off to infinity – what we mean by this will be explained below.

We could similarly ask how interactions modify the photon mass. If one naively introduces the same cut-off to regularize the integrals for vacuum polarization, it is well known that the corrections to the squared mass of the photon are quadratically divergent. This would imply that electromagnetic gauge invariance is broken! The cut-off is, however, not a gauge invariant regulator.[3] If instead we use a gauge-invariant regulator (such as dimensional regularization) the radiative correction to the photon mass vanishes. We say that gauge invariance protects the photon from acquiring a mass.[4] Indeed in quantum electrodynamics of fermions, the leading divergence in *any* quantity is logarithmic, as the reader may readily argue from dimensional analysis. There are no quadratic or linear divergences.

The divergence structure of field theories with elementary scalars is, however, quite different. To see this, we will examine the radiative corrections to the scalar boson mass in the SM. We could equally well have performed the same analysis for scalar electrodynamics defined by (1.1b) and arrived at the same conclusion. Examples of one-loop corrections to m_H are shown in Fig. 2.1. The first of these graphs gives a momentum independent energy shift to the single scalar boson state. Using standard quantum mechanics perturbation theory, this may be evaluated by computing the diagonal element of the interaction Hamiltonian given by the second term in (1.18). Since the answer is independent of the momentum of the state, it corresponds to a mass correction. We find,

$$\delta m^2_{H_{\rm SM}} = \langle H_{\rm SM}| \frac{g^2 m^2_{H_{\rm SM}}}{32 M_W^2} H^4_{\rm SM} |H_{\rm SM}\rangle = 12 \frac{g^2 m^2_{H_{\rm SM}}}{32 M_W^2} \int \frac{{\rm d}^4 k}{(2\pi)^4} \frac{\rm i}{k^2 - m^2_{H_{\rm SM}}},$$

$$= \frac{12 g^2 m^2_{H_{\rm SM}}}{32 M_W^2} \frac{1}{16\pi^2} \left(\Lambda^2 - m^2_{H_{\rm SM}} \ln \frac{\Lambda^2}{m^2_{H_{\rm SM}}} + \mathcal{O}(\frac{1}{\Lambda^2}) \right), \qquad (2.1)$$

[3] This is simple to see. Under gauge transformations, $A_\mu(x) \to A_\mu(x) - \partial_\mu \alpha(x)$, or in k-space, $A_\mu(k) \to A_\mu(k) - {\rm i}k_\mu \alpha(k)$, for all values of k. We see that even if we cut off the modes with $k > \Lambda$ in one gauge, these do not vanish in a different gauge – i.e. a cut-off on the momentum integrals is a gauge-dependent notion.

[4] We are tacitly ignoring the possibility of dynamical gauge symmetry breaking first pointed out by J. Schwinger, *Phys. Rev.* **125**, 397 (1962).

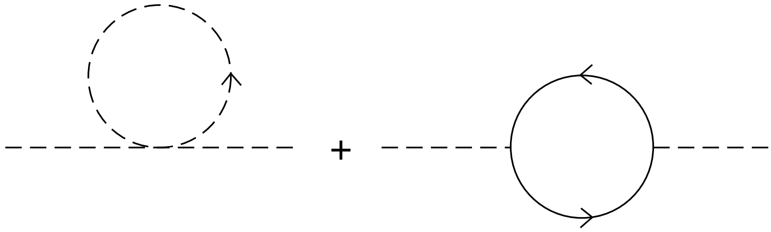

Figure 2.1 Examples of quadratically divergent Feynman graphs contributing to the corrections to the Higgs boson mass in the Standard Model.

where 12 comes from 12 possible field contractions. We see that this correction is quadratically divergent. Mass corrections are more conventionally computed by evaluating the correction to the scalar propagator. We will do so in Chapter 3 for scalar fields in a simple supersymmetric model introduced therein, and see that, once again, there are quadratically divergent contributions to the scalar boson mass.

Within the SM, there are other quadratically divergent contributions to the scalar mass from gauge boson loops, as well as from fermion loops. The computation of the gauge boson loop contribution is very similar to the one that we have just performed, and will be left as an exercise for the reader. We will not evaluate the fermion loop contribution at this point since we perform a very similar calculation in the next chapter. We only note that we expect the top loop contribution to dominate, and further, that this correction is also quadratically divergent, but with opposite sign for the Λ^2 term to the Higgs loop contribution. This difference in signs between contributions from boson and fermion loops is a general feature.

Exercise *Define the chiral projections $\psi_{L/R} \equiv P_{L/R}\psi$, where $P_L = \frac{1-\gamma_5}{2}$ and $P_R = \frac{1+\gamma_5}{2}$. For any two Dirac spinors ψ and χ, verify that,*

$$\bar{\psi}\gamma_\mu\chi = \overline{\psi_L}\gamma_\mu\chi_L + \overline{\psi_R}\gamma_\mu\chi_R,$$
$$\bar{\psi}\chi = \overline{\psi_L}\chi_R + \overline{\psi_R}\chi_L.$$

Convince yourself that these identities imply that the kinetic energy term and the interaction term in (1.1a) preserve chirality, while the mass term does not. Would chirality be conserved had the interaction been axial vector instead of vector?

2.2 Why is the TeV scale special?

Having established the difference between the divergence structure of quantum field theories with and without elementary scalar fields, let us examine the sense

https://doi.org/10.1017/9781009289801.003 Published online by Cambridge University Press

in which this difference may be of relevance. We note at the outset that quadratic divergences do not imply any logical problem. The SM is a renormalizable quantum field theory, and we can use it to evaluate radiative corrections to any precision that we may choose. Indeed, we have already observed that unless these corrections are included, SM predictions are at variance with precision measurements from LEP and other colliders. What then is the problem?

To understand this, we write the one-loop corrected physical Higgs boson mass as,

$$m^2_{H_{\rm SM}}(\text{phys}) \simeq m^2_{H_{\rm SM}} + \frac{c}{16\pi^2}\Lambda^2, \tag{2.2}$$

where $m^2_{H_{\rm SM}}$ is the Higgs mass squared parameter in the Lagrangian and the second term denotes the quadratically divergent correction in (2.1); we have dropped the $\ln\Lambda$ terms in writing this formula. The coefficient c depends on the various coupling constants of the SM. In writing (2.2), we only integrate over the energy–momentum range for which we expect the SM to provide a reasonable description. In other words, we interpret the cut-off Λ as the scale at which the SM ceases to be valid. This may be because new degrees of freedom that are not included in the SM begin to become important. These new degrees of freedom may be unknown heavy particles whose effects are negligible at low energy, for instance, new particles associated with grand unification. It is also possible that the SM breaks down because of new form factors (whose origin may be some unknown strongly coupled dynamics) that develop at the scale Λ. The scale Λ might be as low as several TeV, but certainly no higher than the reduced Planck scale $M_{\rm P} \simeq 2.4 \times 10^{18}$ GeV, the scale at which quantum gravity corrections are expected to become important.

How do we judge what values of Λ are reasonable? *Perturbative unitarity* arguments imply that the physical Higgs boson mass $m_{H_{\rm SM}}(\text{phys})$ that appears on the left-hand side of (2.2) has to be smaller than a few hundred GeV.[5] If we now require (2.2) to be satisfied without excessive fine tuning between the two terms on its right-hand side, we would have to deduce that $\Lambda \leq \mathcal{O}(\text{TeV})$. Again, we stress that this conclusion stems not from a logical inconsistency of the SM but from the additional "no fine-tuning requirement" that we impose on our theory. In the SM the fine tuning that is required can be truly incredible: if we assume the validity of the SM as a low energy effective theory below the GUT scale, and take $\Lambda = M_{\rm GUT} \sim 10^{16}$ GeV, then the Lagrangian mass parameter $m^2_{H_{\rm SM}}$ will have to be fine-tuned to 1 part in 10^{26} to provide the needed cancellation that will maintain a physical Higgs mass below its unitarity limit. In contrast, the logarithmic term in (2.1) contributes a correction which is $\sim m^2_{H_{\rm SM}}$ even for $\Lambda \sim M_{\rm P}$. This is what

[5] D. Dicus and V. Mathur, *Phys. Rev.* **D7**, 3111 (1973); B. Lee, C. Quigg and H. Thacker, *Phys. Rev.* **D16**, 1519 (1977).

we meant when we said above that the $\ln \Lambda$ corrections, which are also present for fermions, are not large. Put somewhat differently, the large Λ^2 corrections imply that if we use the high energy theory (from which the SM originates as the effective low energy theory) to make predictions at TeV energies, these predictions would be extremely sensitive to the parameters of the high energy theory if $\Lambda \gg 1$ TeV. While this is not logically impossible, this is usually thought to be symptomatic of a deeper problem. We refer to this as the *fine-tuning problem* of the SM. If we take this seriously, we are led to the conclusion that *there must be new degrees of freedom that manifest themselves in high energy collisions at the TeV energy scale.* This is especially exciting because we expect this scale to be directly probed by experiments at the Large Hadron Collider (LHC) which is expected to begin operation at CERN in 2007.

2.3 What could the New Physics be?

Although the arguments that we have made point to the existence of new degrees of freedom at the TeV scale, they do not by themselves provide clues as to what this new physics might be. Before proceeding further, however, let us carefully examine if it is possible to evade the conclusion about the existence of new physics at the TeV energy scale.

1. An obvious out is to accept that nature is fine-tuned, and proceed. We would then have to give up on deducing the parameters of the low energy theory from high scale physics, in the manner that we were able to deduce the Fermi coupling G_F in terms of the parameters of electroweak gauge theory. Some authors have recently suggested such a philosophy. They argue that while there have been several proposals that allow us to solve the fine-tuning problem of the SM that we have discussed, it is still necessary to fine-tune the cosmological constant, and to much greater extent than the scalar mass. They argue that accepting the greater fine tuning but not the lesser one seems artificial. We will not pursue this line of reasoning any further here.
2. There are no elementary scalar fields in nature (and hence no associated fine-tuning issues), but composite states of tightly bound fermions play the role of the Higgs boson. This is the idea behind the technicolor model, which posits new *technifermions* that interact and bind together via an asymptotically free QCD-like technicolor interaction, that becomes confining at the TeV scale.[6] Technicolor dynamics causes technifermions to condense, leading to a breakdown of electroweak gauge invariance. This attractive picture runs into problems when we attempt to use the same idea to give masses to fermions.

[6] For reviews, see e.g. E. Fahri and L. Susskind, *Phys. Rep.* **74**, 277 (1981) and K. Lane, hep-ph/0202255.

https://doi.org/10.1017/9781009289801.003 Published online by Cambridge University Press

Construction of realistic models with massive fermions requires the introduction of yet other *extended technicolor* interactions. The simplest models predict flavor-changing neutral current processes at unacceptably large rates. More complicated technicolor theories can be constructed to avoid these problems, but the models become cumbersome and are often in conflict with data involving precision electroweak measurements. We remark that the technicolor approach predicts several pseudo-Goldstone bosons with masses in the several hundred GeV to TeV range as their signature.

3. The arguments leading to (2.2) are inherently perturbative, and would break down if interactions of Higgs bosons with themselves or with the gauge or fermion sector became strong at the scale Λ. One might then expect that these strong interactions would result in new resonances, especially in the Higgs boson and electroweak gauge boson scattering amplitudes, that would reveal themselves in high energy collisions. Even if there is no resonance, just an increase in the scattering amplitude may be experimentally observable in some channels, for instance in $W^{\pm}W^{\pm}$ scattering, where the background is small.[7]
4. A very radical alternative that has received some attention in the last several years is that gravitational effects become strong at an energy scale close to the weak scale. The motivation is to explain the evident difference in the size of gravitational and gauge interactions. Specifically, it is envisioned that there are additional compact spatial dimensions with a size $M_c \ll M_{\rm P}$, and that (unlike SM gauge interactions) gravitational interactions permeate all these dimensions. Gravity appears weak at distances $\gg M_c^{-1}$ because most of the flux is "lost" in these additional dimensions. These extra dimensions are directly probed in particle collisions at energy scales $\gtrsim M_c$ where effects of gravitation become important. In such a scenario, the cut-off Λ in (2.2) would be $\sim M_c$ and the fine-tuning problem disappears if $M_c \sim \mathcal{O}(\text{TeV})$. In this case, we would expect to see exotic effects from Kaluza–Klein resonances of ordinary particles, production of black holes of masses a few times M_c, and strong gravity at high energy colliders.[8]
5. We have implicitly assumed that if a theory has quadratic divergences, these will necessarily show up at the lowest order. It is logically possible that the quadratic divergence only appears at the multi-loop level, but cancels at the one-loop level. In this case, the second term on the right-hand side of (2.2) will have additional powers of $16\pi^2$ in the denominator, and the scale Λ will be correspondingly pushed up. Scenarios where quadratic divergences appear only at the two-loop level have recently been constructed, and go under the rubric

[7] See e.g. J. Bagger *et al.*, *Phys. Rev.* **D49**, 1246 (1994) and *Phys. Rev.* **D52**, 3878 (1995).

[8] For principles and overview, see C. Csaki, hep-ph/0404096; see also J. Hewett and M. Spiropulu, *Ann. Rev. Nucl. Part. Sci.* **52**, 397 (2002) for a review of the phenomenological implications of extra dimensional models.

of *Little Higgs* models.[9] They require new particles to ensure this cancellation. It appears that the parameter space of the simplest of these models is already severely constrained by experimental data. Also, it is clear that this approach only postpones the problem by a couple of orders of magnitude in energy, so that we will need yet more new physics at the 100 TeV scale.

We see that all but the first of these alternatives to attempt to evade the perturbative argument that led us to conclude that there would be new degrees of freedom that could be explored at the TeV scale invoke new strong interactions at this scale, and so lead to potentially observable signatures in particle scattering at $\sqrt{s} \sim \mathcal{O}(\text{TeV})$. Naively, we would also expect that these strong interactions might not decouple, and potentially lead to observable effects in precision measurements at LEP.

The alternative is to assume that the arguments that lead to (2.2) are valid, and that there are new degrees of freedom that are *perturbatively coupled* to the particles of the SM. These new degrees of freedom must then serve to cancel the quadratic divergence that appears because of the presence of elementary scalars. Moreover, it is desirable to have this cancellation occur to all orders, not just at the one-loop level. We have already had a glimmer of how such a cancellation may be arranged when we remarked that the boson and fermion loops led to opposite signs for the coefficient of the Λ^2 term in (2.2). However, in general, one would not expect this cancellation to be complete except by accident, and further, even if we somehow got fermion loops to cancel boson loops at the one-loop level, we would not expect the cancellation to continue at higher orders, *unless the couplings of fermions and bosons are somehow related*. Such relations occur only due to symmetries.[10] However, all symmetries that we know (Lorentz symmetry, baryon or lepton $U(1)$ symmetry, isospin, ...) relate the properties of bosons (fermions) to those of other bosons (fermions), but never those of bosons to those of fermions. A symmetry that relates properties of bosons and fermions is a truly novel symmetry referred to as a *supersymmetry*.

We will see that supersymmetry requires that for every boson, a fermion partner should exist, and vice versa. In other words, every SM boson (fermion) has an as yet unseen fermionic (bosonic) supersymmetric partner. Moreover, supersymmetry relates the interactions of SM particles and their supersymmetric partners in the same way that isospin relates the interactions of protons and neutrons. It is these

[9] For a review, see M. Schmaltz, hep-ph/0210415.

[10] Relativistic quantum electrodynamics is a well-known precedent for a symmetry changing the divergence structure of the theory. This theory has charge conjugation invariance, and requires the presence of anti-particles. In classical theory, it is well known that the self energy of the electron diverges inversely as its size (or equivalently, linearly as the cut-off Λ). The existence of the positron is irrelevant for the classical calculation. The quantum fluctuations include the effects of all particles and thus know about the positron whose existence serves to cancel the linear divergence, leaving the electron self energy that depends logarithmically on the size of the electron.

supersymmetric particles that serve as the new perturbatively coupled degrees of freedom that act to cancel the quadratic divergences of the SM.

We should make a clear distinction between the fine-tuning problem, and the issue of *the origin* of the very small ratio between the weak and GUT or Planck scales. Supersymmetry, by itself, does not explain *why* we have such a small dimensionless ratio. However, once we introduce this ratio by choosing the parameters of the Lagrangian accordingly, this ratio is not destabilized by radiative corrections if the theory is supersymmetric, or if (as we will see in Chapter 3 and again in Chapter 7) supersymmetry is broken appropriately. In the literature, this is also stated by saying that supersymmetry solves the *technical aspect* of the gauge hierarchy problem that is endemic to theories with elementary scalar fields.

Supersymmetry was discovered in the late 1960s and early 1970s under quite different motivations.[11] The first four-dimensional globally supersymmetric quantum field theory was written down in 1974 by Wess and Zumino, and supergravity was discovered shortly thereafter (though attempts to construct locally supersymmetric theories had been made much earlier). While Fayet had already pioneered many phenomenological studies of supersymmetry, it was only after it was recognized that supersymmetry provided a solution to the fine-tuning problem of the SM that there was an explosion of interest in the particle physics community.[12] The simplest viable supersymmetric version of the SM – the Minimal Supersymmetric Standard Model, or MSSM – was developed in the early 1980s, and has since served as the starting point for many phenomenological analyses.

Although it is possible that the fine-tuning problem is solved by something no one has yet thought of, from our vantage point today, the motivations for examining supersymmetry are numerous, and remain as strong as ever.

- *Aesthetics*
  Two pillars upon which the fundamental laws of physics are formulated are the theories of relativity and quantum mechanics. These two are successfully merged within the highly constrained (and hence predictive) framework of relativistic quantum field theory. Haag, Lopuszanski, and Sohnius generalized earlier work by O'Raifeartaigh and by Coleman and Mandula, and showed that the most general symmetries of the S-matrix are a direct product of the super-Poincaré group, which includes supersymmetry transformations linking bosons with fermions in addition to translations, rotations and boosts, with the internal symmetry group. It would be a pity if nature did not make use of this additional mathematical structure at some level.

[11] For an early history, see G. L. Kane and M. Shifman, *The Supersymmetric World: The Beginning of the Theory*, World Scientific (2000).
[12] E. Witten, *Nucl. Phys.* **B188**, 513 (1982); R. Kaul, *Phys. Lett.* **B109**, 19 (1982).

https://doi.org/10.1017/9781009289801.003 Published online by Cambridge University Press

- *Ultra-violet behavior and fine tuning*
  We have already seen that, unlike the case of the Standard Model, the scalar potential of supersymmetric models is stable under radiative corrections, provided that supersymmetric particle masses are comparable to the weak scale. Supersymmetric grand unified models are thus technically natural.
- *Connection to gravity*
  By elevating global supersymmetry transformations to local ones, one is forced into introducing a spin 2 massless gauge field, the graviton, which mediates gravitational interactions (together with its superpartner, the gravitino) in much the same way as local gauge invariance requires us to introduce gauge bosons. Just as gauge invariance is sufficient to fix the dynamics, local (super)symmetry dictates the dynamics of supergravity, which includes Einstein's general relativity. Like any four-dimensional theory of gravity, this supergravity theory is not renormalizable. Nonetheless, this connection to gravity is very tantalizing.
- *Ultra-violet completeness*
  Except for supersymmetry, all the extensions of the SM that we considered above to ameliorate the fine-tuning problem are effective theories and require even more new physics at a scale that is only a couple of orders of magnitude above the TeV scale.[13] While this is not necessarily an argument against such scenarios, the fact that supersymmetric theories can in principle be extrapolated all the way to the GUT or Planck scales is especially attractive.
- *Connection to superstrings*
  Supersymmetry is an essential ingredient of superstring theories, thought to be candidates for a consistent, finite quantum theory of gravitation. In this framework, the problem of non-renormalizability of gravitational theory is bypassed by moving away from point-like particles to intrinsically finite theories of extended objects, open or closed loops of a fundamental string. Superstring theories have enjoyed an important success in the counting of microscopic states of a string black hole.

In addition to these aesthetic considerations, there are several experimental arguments that also highlight the promise of supersymmetric models.

- *Unification of gauge couplings*
  The values of running gauge couplings measured at LEP do not unify if we evolve these to high energies using the renormalization group equations of the Standard Model. If instead, we extrapolate the gauge couplings from $Q = M_Z$ to $Q = M_{\rm GUT}$ using supersymmetric evolution equations, they unify remarkably well provided superpartner masses are in the range 100 GeV–10 TeV. It is extremely

[13] Models with warped extra dimensions may be a possible exception to this.

https://doi.org/10.1017/9781009289801.003 Published online by Cambridge University Press

suggestive that this scale agrees so well with the scale that we inferred from fine-tuning arguments. Unless this is a perverse accident, this strongly points toward supersymmetric Grand Unification. Moreover, the value of M_{GUT} that is obtained is somewhat higher than in non-supersymmetric GUTs; this reduces the amplitude for proton decay by GUT boson exchange, bringing this contribution in accord with lower limits on the proton lifetime. Potentially larger superpotential contributions to the proton decay amplitude have to be controlled, however.

- *Cold dark matter*
  All supersymmetric models with a conserved R-parity quantum number include a stable massive particle which is usually electrically and color neutral, and so makes an excellent candidate for the observed cold dark matter in the Universe.
- *Radiative breakdown of electroweak symmetry*
  In the SM, electroweak symmetry breaking (EWSB) can be accommodated by appropriate choice of the parameters of the scalar potential, without any explanation for this choice. We will see in Chapter 11 that in many supersymmetric models, scalars with the same gauge quantum numbers have the same mass parameters, renormalized at some high scale. Over a large part of the model parameter space, renormalization effects drive the Higgs boson squared mass parameters to negative values, while those for scalars with non-trivial $SU(3)_{\mathrm{C}} \times U(1)_{\mathrm{em}}$ are left positive, resulting in the observed electroweak symmetry breaking pattern. Radiative electroweak symmetry breaking, as this mechanism is known, occurs naturally if $m_t \sim 100$–200 GeV. While this mechanism is very attractive, it cannot be regarded as a complete explanation since there are parameter ranges for which color and electromagnetic gauge invariance may be broken. We should add that it also requires that the soft SUSY breaking parameters are $\sim$TeV.
- *Decoupling in SUSY theories*
  Radiative corrections from SUSY particles in loops to electroweak observables in LEP experiments rapidly decouple with SUSY particle masses, so that SUSY models can readily replicate the apparent successes of the SM in explaining the LEP data. This is not to say that *all* SUSY loops decouple. For instance, SUSY loop corrections to the couplings of quarks and leptons to Higgs scalars (and hence to the relation between fermion masses and Yukawa couplings) do not necessarily decouple.
- *Mass of the Higgs boson*
  The Higgs boson of the SM can have mass of any value between the lower limits set by LEP and LEP2 experiments ($m_{H_{\mathrm{SM}}} > 114.4$ GeV), up to ~ 800 GeV. The much more constrained Higgs sector of the MSSM requires the lightest Higgs scalar h to have mass $m_h \lesssim 135$ GeV. Meanwhile, precision measurements of electroweak parameters which are sensitive to the mass of the Higgs boson point towards $m_h \sim 120$ GeV, with $m_h \lesssim 200$ GeV at the 95% CL.

https://doi.org/10.1017/9781009289801.003 Published online by Cambridge University Press

The wide array of issues that are addressed by the inclusion of supersymmetry in particle physics has led many physicists to suspect that supersymmetry is realized in nature, perhaps at or just above the TeV energy scale. The goal of this book is to show how weak scale supersymmetry can be incorporated into the basic laws of physics, and detail how to extract the observable consequences of the supersymmetry hypothesis. The important and exciting conclusion is that the idea of weak scale supersymmetry can be tested at various collider and non-accelerator experiments.