Naturalness Under Stress

Michael Dine

Santa Cruz Institute for Particle Physics and Department of Physics, University of California, Santa Cruz, California 95064

Annu. Rev. Nucl. Part. Sci. 2015. 65:43–62

First published online as a Review in Advance on May 4, 2015

The *Annual Review of Nuclear and Particle Science* is online at nucl.annualreviews.org

This article's doi:
10.1146/annurev-nucl-102014-022053

Keywords

naturalness, supersymmetry, technicolor, beyond the Standard Model, fine-tuning

Abstract

Naturalness has for many years been a guiding principle in the search for physics beyond the Standard Model, particularly for understanding the physics of electroweak symmetry breaking. However, the discovery of the Higgs particle at 125 GeV, accompanied by the exclusion of many types of new physics expected in natural models, has called the principle into question. In addition, apart from the scale of weak interactions, there are other quantities in nature that appear unnaturally small and for which we have no proposal for a natural explanation. I first review the principle, then discuss some of the conjectures it has spawned. I then turn to some of the challenges to the naturalness idea and consider alternatives.

Contents

1. NATURALNESS: A CONTEMPORARY IMPLEMENTATION OF DIMENSIONAL ANALYSIS

In our first science courses, we learn about the importance of dimensional analysis. Often this is presented as a consistency requirement for calculations of physical quantities, but it shapes our understanding of physical systems in a fundamental way. For example, from the electron mass, m_e, the speed of light, c, and Planck's constant, $\hbar$, we can form a quantity with dimensions of length $a = \frac{\hbar}{mc} \approx 4\times 10^{-15}$ cm, supplemented with the insight that the strength of the force between the proton and the electron is proportional to e^2, so the size should get larger as e^2, or $a = \frac{\hbar}{mce^2}$. To know the exact coefficient—which is an order-one number—we need to solve the Schrödinger equation completely. But we get a nice qualitative, and rough quantitative, picture without much trouble.

Similarly, the size of atomic nuclei is large compared with the Compton wavelength of the proton, suggesting that there should be physics associated with this larger length scale. Even if we are ignorant of the detailed mechanism, this suggests the existence of a particle with a mass roughly equal to that of the pion.

This sort of reasoning has had successes in many other areas of physics. More interesting are questions for which it fails, at least at first sight. In 1899, Planck (1) noted that from $\hbar$, c, and $G_{\rm N}$ (the Newton constant), one can form a quantity with units of mass $M_p = \sqrt{c\hbar/G_{\rm N}}$. At or below this scale, quantum mechanical effects should be important in general relativity.

Suppose that there is some underlying theory, which includes general relativity, from which one computes the electron mass. Dimensional analysis would say that $m_e = \beta M_p$ where β is an order-one number. Of course, this is terribly wrong—dimensional analysis fails stupendously here.

Lorentz (2) encountered this issue in a somewhat different way, which provides a different—and equally useful—perspective on the problem. He modeled the electron as a smooth charge distribution with a characteristic size, a. One would expect that the mass of the electron would then be at least of order the self-energy of the electron arising from its Coulomb field, $m_e \approx \frac{e^2}{4\pi a}$. This statement might be viewed as a prediction of a : $a \approx 10^{-10}$ cm or even 10^{-12} cm. But from present-day experiments, we know that $a < 10^{-17}$ cm. At first sight, therefore, this is a serious failure of dimensional analysis. Alternatively, we might describe this issue as a problem of "naturalness" or fine-tuning. If there is an additional, "bare" mass parameter, $m_e^{(0)}$, then $m_e = m_e^{(0)} + \frac{e^2}{a}$. Each term separately is approximately 5×10^4 the observed mass of the electron.

The resolution of this puzzle has been known since the work of Weisskopf (3, 4) in 1934. His supervisor at the time, Wolfgang Pauli, assigned him the problem of computing the corrections to the energy of a free electron due to its interactions with its own fields. Using the newly discovered rules of (relativistic) quantum mechanics, this task required including not only the interaction of the electron with its Coulomb field but also contributions to the energy from intermediate states of two types: one with an electron and a photon and one with two electrons, a positron, and a virtual photon. The expressions were divergent at high energies (corresponding, in modern language, to high virtual photon momenta), and Weisskopf assumed that these were cut off by the size of the electron. In his first attempt, he encountered a similar linear divergence ($1/a$) as in Lorentz's calculation, but following an observation by Furry, he quickly corrected a mistake and found that the leading linear divergence cancelled, leaving only a logarithmic dependence on the cutoff. The full expression, which can be derived by a modern field theory student in a matter of minutes, is

$$m_e = m_e^{(0)}\left(1 - \frac{6\alpha}{4\pi}\log(m_e a)\right). \qquad 1.$$

Even for extremely small a ($a = 10^{-31}$ cm), the correction is only $\sim$20% of the leading result. It is remarkable that the naturalness problem of the classical theory is resolved not simply by the quantization of the theory but by the fact that there are additional degrees of freedom required by the relativistic quantum theory. In fact, if the electron had been a scalar, this would not have happened, as discussed further below; instead, the mass squared diverges quadratically with the cutoff.

It is crucial that Equation 1 be proportional to the original electron mass, the parameter that appears in the Lagrangian for the theory. This proportionality can be understood in a conceptual way. In the limit that the mass of the electron vanishes, quantum electrodynamics (QED) is more symmetric. Setting the mass term, $m_e^{(0)}$, to zero in the usual Dirac Lagrangian

$$\mathcal{L} = \bar{\psi}[i\gamma^\mu(\partial_\mu - ieA_\mu) - m_e^{(0)}]\psi, \qquad 2.$$

one has a symmetry under the chiral or Weyl transformation: $\psi \to e^{i\omega\gamma_5}\psi$. In this limit, all effects in the theory—and, in particular, any corrections to the Lagrangian—must respect the symmetry. This means, in particular, that any correction to the mass must vanish as the mass tends to zero, precisely the feature of Equation 1.

So we see that, although even now we do not have a compelling microscopic explanation of this mass, a small electron mass is special in that QED becomes more symmetric. 't Hooft (5) elevated this to a principle of naturalness: A quantity in nature should be small only if the underlying theory becomes more symmetric as that quantity tends to zero. There are other instances in which this reasoning works remarkably well. Consider, for example, the mass of the proton. The proton is composed of quarks and gluons, but its mass has very little to do with the masses of the quarks, which are of order the small difference between the proton and neutron masses. So, again, we might ask why the mass of the proton is not M_p. The answer turns out, again, to be related to symmetries. Setting the quark masses to zero, we find that the classical action of QCD has no scale—the theory has a symmetry called scale or conformal invariance. If this symmetry were exact, the proton would necessarily be massless; in the quantum theory, this symmetry is broken by a small amount.

The violations of scale invariance are associated with the phenomenon of renormalization in quantum field theory. Renormalization is the statement that the parameters of a theory vary with length or energy scale. This variation is logarithmic, encoded in renormalization group equations. For the strong coupling, α_s, specifically:

$$\frac{\mathrm{d}\alpha_s}{\mathrm{d}t} = -2b_0\alpha_s^2. \tag{3}$$

Here, $t = \log(M/E)$, where M is a UV cutoff (or matching scale) and b_0 is a constant. So, if one asks at what scale $E \equiv \Lambda$, the coupling becomes of $\mathcal{O}(1)$:

$$\Lambda = M_p e^{-\frac{2\pi}{b_0\alpha_s(M_p)}}. \tag{4}$$

For QCD, b_0 is a number of order seven, so if g_s at M_p is $\sim$0.5, the exponential is extremely small, and the scale Λ is of order the proton mass.

Most of the parameters of the Standard Model (SM) are natural in the sense of 't Hooft, but there are some quantities that are not. It is precisely the failures of dimensional analysis that are most interesting. As for the electron and proton masses, they have the potential to point to possible new phenomena in nature—new degrees of freedom, interactions, and/or symmetries. For a long time, these sorts of puzzles have guided speculation about physics beyond the SM.

2. NATURALNESS PROBLEMS IN PARTICLE PHYSICS

Our current theories of the laws of nature are best viewed as tentative effective field theories, valid at energies below some scale at which new degrees of freedom or other phenomena might manifest themselves. Naturalness, from this perspective, is the assertion that features of this effective field theory should not be extremely sensitive to the structure of the underlying theory. For the electron, this is the statement that its Yukawa coupling to the Higgs boson receives only small corrections as one studies the theory at progressively higher energy scales. For the strong interactions, as the existence of a proton much lighter than the Planck scale can be explained by an order-one pure number at M_p.

The masses of the quarks and leptons are controlled by symmetries analogous to those controlling the mass of the electron in the Weisskopf calculation. The $SU(2) \times U(1)$ symmetry of the SM forbids masses larger than $y \times 250$ GeV. Here, y is a pure number, the Yukawa coupling of the quark or lepton. For the quarks and charged leptons, this number ranges from $\sim$1 for the top quark to 10^{-5} for the electron. The spread in these numbers raises many puzzles, but it is not unnatural. Just as in the case of the small electron mass in pure QED, in the limit of very small electron Yukawa coupling the theory becomes more symmetric. Indeed, if we set all of the Yukawa couplings to zero, the theory possesses a large symmetry. Several theories that might account for these small numbers and the hierarchies among them have been proposed. It is fair to say that none is completely compelling by itself, nor does any make unequivocal predictions for experiment. Still, the existence of a hierarchy in fermion masses and mixings does not pose a fundamental conceptual problem.

There is one quantity in the SM that fails 't Hooft's test and raises precisely the sorts of issues posed by the classical theory of the electron. This is the mass of the Higgs particle, which is tied to the scale at which the symmetry of the electroweak theory is broken. In the simplest version of

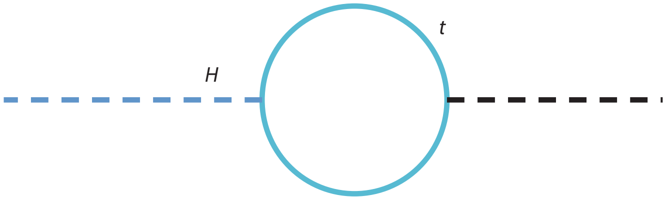

Figure 1

One-loop correction to Higgs mass involving top quarks.

the SM, the potential of the Higgs field is

$$V(\phi) = -\mu^2|\phi|^2 + \frac{\lambda}{4}|\phi|^4. \qquad 5.$$

Assuming that this potential describes the recently observed Higgs particle (and measurements to date are consistent with this picture), we know the values of μ and λ : $\mu \approx 89\,\text{GeV}$, and $\lambda \approx 0.13$.

Dimensional analysis, however, predicts $\mu^2 \approx M_p^2$, and there is no enhancement of the symmetry of the theory if we take $\mu^2 \to 0$. If we repeat Weisskopf's calculation for this case, we confront this issue directly. The strongest coupling of the Higgs field in the SM is its Yukawa coupling to the top quark: $\mathcal{L}_{\bar{t}tH} = y_t H Q_3 \bar{t}$, where Q_3 refers to the third quark doublet, consisting principally of the top and bottom quarks. At one loop, there is a correction to the Higgs mass that comes from the diagram shown in **Figure 1**. This correction is given by

$$\delta\mu^2 = -6y_t^2 \int \frac{d^4k}{(2\pi)^4}\frac{1}{k^2}, \qquad 6.$$

where the integral is an ordinary Euclidean integral, which diverges quadratically. If the cutoff is the Planck scale, then this correction is enormous—consistent with expectations from dimensional analysis, it is ~34 orders of magnitude larger than M_H^2, which corresponds to a fine-tuning of the bare parameters against the radiative correction at the part in 10^{34} level.

This is one of many such contributions to the Higgs mass, which include those from diagrams involving lighter quarks, gauge bosons, and the Higgs bosons themselves. Wilson (6) was the first to raise the issue of the quadratic growth (divergence) in corrections to scalar masses.

3. OTHER "UNNATURAL" STANDARD MODEL PARAMETERS

As mentioned above, the small quark and lepton masses (Yukawa couplings) are natural in the sense of 't Hooft. In the limit that all of the quark and lepton masses vanish, the SM has a large global symmetry. For each type of quark or lepton (where type is defined the gauge quantum number of the associated field), namely $Q_f, \overline{U}_f, \bar{d}_f, L_f$, and $\bar{e}_f$, the theory has a separate $U(3)$ symmetry. This symmetry is defined, in the case of Q_f, for example, by $Q_f \to U_{f,f'} Q'_f$. As a result, quantum corrections to the Yukawa couplings (and, hence, masses) vanish in the limit that the masses tend to zero.[1] Many physicists have explored the possibility that some underlying theory possesses precisely these symmetries (or perhaps a continuous or discrete subgroup), and that they are spontaneously broken by a small amount.

[1]There is an exception associated with the fact that one linear combination of the $U(1)$ subgroups of these symmetries has a QCD anomaly. However, Feynman diagram corrections still vanish, and the resulting effects are quite small.

One small parameter does appear, on its face, to violate 't Hooft's condition. It is possible to add to the QCD Lagrangian the term

$$\mathcal{L}_\theta = \frac{\theta}{16\pi^2} G_{\mu\nu}\tilde{G}_{\mu\nu}, \tag{7}$$

where $G_{\mu\nu}$ is the QCD field strength and $\tilde{G}_{\mu\nu}$ is its dual: $\tilde{G}_{\mu\nu} = \frac{1}{2}\epsilon_{\mu\nu\rho\sigma}G^{\rho\sigma}$. This term is odd under parity (*P*) and even under charge conjugation (*C*), so it violates *CP*. In electrodynamics, the analogous term is $\mathbf{E}\cdot\mathbf{B}$, which is a total derivative and has no effect.[2] In QCD, this term is also a total derivative. As a result, it does not affect the equations of motion. However, it does have physical effects. Most notably, using current algebra one can compute the electric dipole moment of the neutron, d_n, as a function of θ (8):

$$d_n = 5.2 \times 10^{-16}\theta \text{ cm}. \tag{8}$$

From the experimental limit, $d_n < 3 \times 10^{-26}\, e$ cm, one obtains $\theta < 10^{-10}$. If nature respected *CP* in the absence of θ, this small value of a dimensionless number would be natural in the sense of 't Hooft. But nature violates *CP*; indeed, the phase appearing in the Cabibbo–Kobayashi–Maskawa (CKM) matrix is of order one. So, like the Higgs mass, this number requires an explanation.

4. PROPOSED SOLUTIONS TO THE PROBLEM OF THE HIGGS MASS

Over the years, several solutions of the hierarchy problem have been proposed.

4.1. Technicolor

Weinberg (9) and Susskind (10) put forward the first solution to the problem of naturalness of the Higgs mass, closely paralleling the understanding of the hierarchy between the proton mass and the Planck scale. They argued that if the Higgs boson was a composite of fermions, with a binding scale of order 1 TeV, this solution would solve the problem. Susskind dubbed this solution technicolor.

Consider the SM without the Higgs particle and with only a single generation of quarks and leptons—that is, with fermions:

$$Q = \begin{pmatrix} u \\ d \end{pmatrix}; \quad \bar{u} \quad \bar{d}; \quad L = \begin{pmatrix} \nu \\ e \end{pmatrix}; \quad \bar{e}. \tag{9}$$

Neglecting the weak coupling, the quark sector of the theory possesses a global symmetry: $SU(2)_\mathrm{L} \times SU(2)_\mathrm{R} \times U(1) \times U(1)$. $SU(2)_\mathrm{L}$ is simply the $SU(2)$ of weak interactions, which rotates the doublet Q; $SU(2)_\mathrm{R}$ is an approximate symmetry under which $\bar{u}$ and $\bar{d}$ transform as a doublet. The $U(1)$ of the SM is a combination of the diagonal generator of the $SU(2)_\mathrm{R}$ and one of these $U(1)$s. The strong interactions break the symmetry into the diagonal subgroup, the familiar $SU(2)$ of isospin, and a $U(1)$; a linear combination of these symmetries is electric charge.

Because the $SU(2)_\mathrm{L} \times U(1)$ subgroup of this symmetry is gauged, the W and Z bosons gain mass, and the photon remains massless. This situation is nicely illustrated with the familiar nonlinear Lagrangian description of chiral symmetry breaking, in which the pions are described by a matrix

[2] This is not quite true; if there are magnetic monopoles in nature, a θ_{QED} parameter affects the properties of their charged excitations, the so-called dyons (7).

of fields with a simple transformation property under the $SU(2)_L \times SU(2)_R$.:

$$\Sigma = e^{i\frac{\pi^a \sigma^a}{2 f_\pi}}; \quad \Sigma \to U_L \Sigma U_R. \qquad 10.$$

The Lagrangian for Σ is

$$\mathcal{L}_\Sigma = f_\pi^2 \mathrm{Tr}(D_\mu \Sigma D_\mu \Sigma^\dagger). \qquad 11.$$

It is instructive to work out the form of the covariant derivatives [the reader for whom this is not familiar would do well to first do the exercise of simply gauging the $SU(2)_L$, where the gauge interactions act only from the left, then include the $U(1)$ by gauging a subgroup of the $SU(2)_R$]. By doing so, one immediately finds that the gauge boson masses are simply those of the SM, with the Higgs expectation value, v, replaced by f_π.

The technicolor hypothesis simply replaces the ordinary quarks by techniquarks, and color by a new interaction: $f_\pi \to F_{tc} = v$. This theory solves the hierarchy problem both in the sense that there are no longer quadratic divergences (loosely, the divergences are cut off at the technicolor scale) and in the sense that it provides an explanation of the weak scale, analogous to the QCD explanation of the proton mass: $F_{tc} = M e^{-\frac{8\pi^2}{b_{tc} g_{tc}(M)^2}}$.

Although this proposal is a beautiful idea, it runs into a number of difficulties. For instance, in this simple form, it has no mechanism to account for the masses of quarks and leptons. One can try to resolve this problem by introducing further gauge interactions, whose role is to break the chiral symmetries that protect fermion masses. The resulting models are quite baroque, requiring many gauge groups and intricate dynamics, but aesthetic objections aside, they run into serious issues with flavor-changing neutral-current processes. Put simply, the SM possesses a variety of approximate symmetries due to small quark masses, and these account, for example, for the small rate for $K \leftrightarrow \bar{K}$ mixing; it is difficult to mimic this phenomenon in a strongly interacting theory.

Prior to the discovery of the Higgs boson, other serious problems had long been noted, especially difficulties with precision studies of the SM (11). The existence of a Higgs boson much lighter than 1 TeV, and with a width less than a few GeV, is particularly difficult to understand in a technicolor framework. Most proposals to explain the existence of a Higgs boson in this framework assume that the technicolor theory is nearly conformal over a range of scales, with a light, SM-like Higgs boson a consequence.

4.2. Little Higgs and Similar Models

An approach that attempts to reconcile the idea of dynamical electroweak symmetry breaking with the existence of a Higgs particle that is light compared with the scale of the new interactions is to consider the Higgs boson an approximate Goldstone boson (12–14). The basic idea of such little Higgs models is that there are some new strong interactions, at a scale M, and that these interactions possess an approximate global symmetry that is spontaneously broken. One of the Goldstone bosons of this symmetry acts as the Higgs boson. The SM gauge interactions necessarily break these symmetries and give rise to a potential. The Higgs mass term induced is too large unless one introduces additional features in such a way that the approximate symmetries are violated only by combinations of additional gauge symmetries. Accounting for fermion masses and satisfying other constraints are challenging.

4.3. Large Extra Dimensions

In the large extra-dimension models (15, 16), one alters the nature of the hierarchy problem by postulating that the fundamental scale of physics is near the scale of electroweak breaking, of order 1 TeV. This proposal can be accommodated if one supposes that there is some number, d, of compact extra dimensions of space (minimally two), with volume ℓ^d. Then, starting with the $(d+4)$-dimensional Einstein action

$$\mathcal{L}_{d+4} = \kappa_{d+4}^{-2} \int d^4x d^d y \sqrt{g} R, \tag{12.}$$

where κ^2 is the $(d+4)$-dimensional Newton constant and y are the extra dimensions, the four-dimensional Newton constant is simply $G_{\rm N} = \kappa^2/\ell^d$. If $\kappa^2 = (\text{TeV})^{-(2+d)}$, then for $d = 2$, for example, the dimensions are of order millimeters; for $d = 6$, one has $\ell \approx 0.2\ \text{MeV}^{-1} \approx 10^{-9}$ cm.

So as not to have a similar dilution of the strength of the SM interactions by ℓ^d, these theories need an additional feature: The SM must exist on a geometric object known as a 3-brane. P-branes are generalizations of membranes (2-branes), strings (1-branes), and particles (0-branes). These 3-branes fill all of space; excitations on the 3-brane behave like particles in four dimensions. The SM gauge bosons, fermions, and Higgs boson, in this picture, are excitations of the brane.

These models make exciting predictions (17). In the two-dimensional case, for example, one predicts a modification of Newton's laws at millimeter scales, which has prompted experimental searches that have constrained the possibility by verifying Newton's laws to remarkably small distances (18, 19). Such models also predict the existence of many new particles, associated with the modes of the higher-dimensional fields on the compact volume (Kaluza–Klein modes). At sufficiently high energies, one should produce large numbers of these particles, essentially uncovering the physics of the higher-dimensional space-time.

This approach alters the question of hierarchy to the question of why these extra dimensions are so large. We have yet to develop a compelling picture, but the possibility is intriguing, and possible short-distance modifications of general relativity or signals of large extra dimensions in accelerators remain active subjects of investigation.

4.4. Warped Extra Dimensions

Warped extra dimensions incorporate elements of the large extra-dimension picture and of technicolor models (20, 21). Here, one has extra dimensions (for simplicity, we consider one extra dimension) and 3-branes. In a simple version, the SM sits on one of two branes. Under certain conditions, the Einstein equations in the higher-dimensional space admit a solution in which the metric varies exponentially with the distance from one or the other brane. This variation is analogous to the variation of couplings with scale in non-Abelian gauge theories (discussed above). The strength of gravity relative to the weak interactions is exponential in the separation of the branes, $e^{-\ell}$. As a result, gravity is very weakly coupled on one brane and strongly coupled on the other. Variants of this idea have some of the SM fields in the bulk space between the branes. Scenarios exist that would account for quark and lepton masses and the suppression of flavor-changing processes. Precision electroweak physics and the observed Higgs particle pose significant challenges, as does embedding this picture in a more complete theory such as string theory. Signals of such warped dimensions include low-lying Kaluza–Klein states, and a great deal of effort has gone into searching for such particles.

5. SUPERSYMMETRY

In implementing 't Hooft's notion of naturalness, we have so far considered symmetries of a sort familiar from quantum mechanics, generated by a charge operator that is a scalar under rotations. But there is another type of symmetry, allowed by general principles of quantum mechanics and relativity, in which the symmetry generators are spinors. This symmetry is known as supersymmetry. We consider it first as a global symmetry, but the symmetry can be elevated to a local, gauge symmetry.

Supersymmetry has many remarkable properties. One is the algebra of the symmetry generators; these obey anticommutation relations with the energy and momentum on the right-hand side:

$$\{Q_\alpha, Q^*_{\dot\beta}\} = P_\mu \sigma^\mu_{\alpha\dot\beta}. \qquad 13.$$

Here, P_μ is the total four-momentum of the system. We are using two-component spinor notation, where $\sigma^i_{\alpha\dot\beta}$ are the ordinary Pauli matrices and σ^0 is the identity matrix. Taking the trace of both sides gives

$$\sum_\alpha Q_\alpha Q^*_\alpha + Q^*_\alpha Q_\alpha = 2E. \qquad 14.$$

As for any symmetry, these generators (charges) commute with the Hamiltonian. Acting on bosonic or fermionic states yields the following relations:

$$Q_\alpha |B\rangle \propto \sqrt{E}|F\rangle; \quad Q_\alpha |F\rangle \propto \sqrt{E}|B\rangle. \qquad 15.$$

As a result, if the symmetry is exact and unbroken, fermions and bosons are degenerate.

This feature of supersymmetry makes it particularly interesting for the hierarchy problem. A fundamental Higgs scalar provides a very simple way to understand quark and lepton masses; it has the additional advantage of being consistent with precision electroweak studies and now with the discovery of what appears to be an elementary Higgs scalar. So, it would be desirable to find theories in which the masses of elementary scalar fields are protected by symmetries. Supersymmetry is the only known such symmetry. As explained above, it is natural for fermions to be light; in the presence of supersymmetry, it follows that it is also natural for bosons, and in particular scalars, to be light.

Of course, there is no such degeneracy among the particles of the SM, so the symmetry must be broken. To account for the Higgs mass, in the spirit of 't Hooft's principle, the breaking scale should be much smaller than the Planck scale. Witten (22) pointed out in 1981 that supersymmetry is particularly susceptible to small, spontaneous breaking. Equation 14 shows that supersymmetry is unbroken if and only if $E = 0$. It turns out that supersymmetric field theories for which $E = 0$ classically have $E = 0$ (and unbroken supersymmetry) to all orders in perturbation theory (23). But this need not hold beyond perturbation theory, and often does not. This means that the energy scale of supersymmetry breaking can take the form $E = Me^{-\frac{8\pi^2}{g^2}}$, which is reminiscent of other hierarchies we have encountered. This phenomenon is referred to as dynamical supersymmetry breaking.

5.1. Basics of Supersymmetric Field Theories

Many excellent texts and review articles on supersymmetry have been published. There is not enough space here to fully elucidate the structure of supersymmetric theories, but a few basic features will be helpful for the subsequent discussion.

1. Supersymmetry multiplets. In globally supersymmetric models, there are two basic types of multiplets: chiral multiplets, consisting of a complex scalar and a spin-1/2 fermion, and vector multiplets, consisting of a chiral fermion and a gauge boson.
2. Interactions between the matter fields. These interactions are described by a holomorphic (analytic) function of the chiral fields (scalar components) called the superpotential, $W(\phi_i)$.

In terms of W, there is a contribution to the potential for the scalars:

$$V_W = \left| \frac{\partial W}{\partial \phi_i} \right|^2, \qquad 16.$$

as well as mass terms and Yukawa couplings for the fermions:

$$\mathcal{L}_f = \frac{1}{2} \frac{\partial^2 W}{\partial \phi_i \partial \phi_j} \psi_i \psi_j + \mathrm{cc}, \qquad 17.$$

where the ψ_i are the fermionic partners of the ϕ_i and cc refers to the complex conjugate.

3. Interactions between the gauge fields. In a particular gauge (the Wess–Zumino gauge), the vector fields interact with one another just as in ordinary non-Abelian gauge theories; the gauginos, λ^a, couple to the gauge fields, as expected for fermions in the adjoint representation.
4. Interactions between the matter fields and the gauge fields. In the same gauge, the scalars and fermions in the chiral multiplets couple to gauge fields just as in ordinary gauge theories. They possess Yukawa couplings to the gauginos:

$$\mathcal{L}_\lambda = \sqrt{2} g (\lambda^a \phi_i^* T^a \psi_i) + \mathrm{cc}. \qquad 18.$$

5. Quartic couplings of scalars charged under the gauge groups:

$$V = \frac{g^2}{2} (\phi_i^* T^a \phi_i)^2. \qquad 19.$$

Supersymmetry can be elevated to a local symmetry. In that case, the gauge field associated with local supersymmetry transformations is the gravitino, $\psi^\mu_\alpha(x)$, a field of spin 3/2. The action becomes distinctly more complicated (24, 25). In the limit of unbroken supersymmetry in flat space, one can define global supercharges, just as one can define a global energy and momentum. These global supercharges still obey the basic algebra (Equation 13). In addition to the chiral and vector multiplets, there is a gravitational multiplet, consisting of the graviton and the gravitino. Small breaking of supersymmetry in supergravity leads to theories that, at low energies, resemble globally supersymmetric theories with explicit soft breaking (26).

5.2. Building Models for Supersymmetry and Its Breaking

If nature is supersymmetric, then the partners of the known fermions (quarks and leptons) are complex scalar fields (with the same gauge charges). These particles are referred to as squarks and sleptons. The partners of the gauge bosons are the gauginos. The fermionic partners of the Higgs fields (supersymmetry requires a minimum of two Higgs doublets) are known as higgsinos.

Constructing realistic models with dynamical supersymmetry breaking poses challenges, so most searches for supersymmetry, and many investigations of the basic features of such theories, start by introducing an explicit, soft breaking of the symmetry. This amounts to simply adding masses for the squarks, sleptons, and gauginos, as well as certain dimensionful couplings (27). These parameters (along with cubic couplings of the scalars) are described as soft because they do not spoil the good UV properties of the theories (26).

In addition to the top quark loop (discussed above), there is now a loop (**Figure 2**) containing a stop that tames the quadratic divergence of the SM. There are actually two types of stops, one from the electroweak doublet and one from the singlet. For simplicity, if we call the mass of each

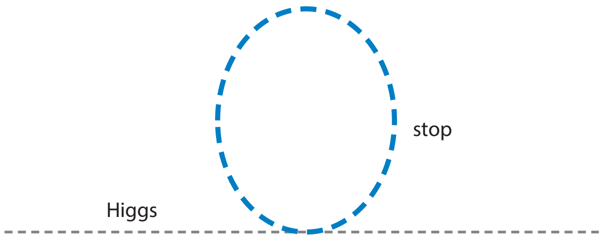

Figure 2

Additional correction to the Higgs mass from stops.

of these scalars $\tilde{m}_t^2$, the two Feynman diagrams yield

$$\delta m_H^2 = 3y_t^2 \int \frac{d^4k}{(2\pi)^4} \left(-\frac{1}{k^2 + m_t^2} + \frac{1}{k^2 + \tilde{m}_t^2} \right). \qquad 20.$$

The minus sign in the first term is the usual minus sign in field theory associated with fermion loops. The leading quadratic divergence cancels, leaving only a logarithmically divergent term:

$$\delta m_H^2 = -\frac{3y_t^2}{16\pi^2} \tilde{m}_t^2 \log(\Lambda^2/\tilde{m}_t^2). \qquad 21.$$

Here, Λ is a UV cutoff, and we have assumed $m_t^2 \ll \tilde{m}_t^2$, consistent with exclusions from the Large Hadron Collider (LHC), discussed in the following section. This situation is closely parallel to that for the electron mass in QED.

6. THE SIMPLEST IMPLEMENTATION OF SUPERSYMMETRY: THE MINIMAL SUPERSYMMETRIC STANDARD MODEL

To develop a supersymmetric phenomenology, we can promote each fermion of the SM to a chiral multiplet and each gauge boson to a vector multiplet. We have quark doublets and antiquark singlets (Q_f, $\overline{u}_f$, and $\overline{d}_f$) and lepton doublets and singlets (L_f and $\bar{E}_f$), where f is a flavor label. There are necessarily two Higgs doublets, H_U and H_D (otherwise, the model is inconsistent), and gauginos accompany each of the gauge bosons.

The superpotential of the model includes couplings of the Higgs particle to quarks and leptons:

$$W = y_{fg}^D Q_f \overline{D}_g H_D + y_{fg}^U Q_f \overline{U}_g H_U + y_f^L L_{\overline{f}} E_f H_D. \qquad 22.$$

The expectation value of the Higgs particle accounts for the fermion masses. There are also various renormalizable terms that can lead to processes in which baryon and/or lepton number is violated. Terms in the superpotential, such as $QL\bar{d}$ and $\bar{u}\bar{d}\bar{d}$, lead to violation of baryon and lepton numbers. If the dimensionless coefficients are of order one, one would expect the proton to decay in $\sim 10^{-24}$ s. Postulating a discrete symmetry, called R parity, forbids these operators. Under this symmetry, all of the particles of the SM, as well as the additional Higgs doublets, are even, whereas all of their superpartners are odd. With this restriction, the minimal supersymmetric Standard Model (MSSM) contains 105 new parameters, associated with the soft breaking of supersymmetry and the additional Higgs field. Consistent with 't Hooft's principle, the R parity–violating couplings might be nonzero but extremely small, leading to a distinctly different phenomenology.

Assuming R parity conservation, the lightest supersymmetric particle is stable. In this case, it must be electrically neutral, presumably some linear combination of the neutral higgsinos and gauginos. The existence of this stable particle implies that production of supersymmetric partners in accelerators is associated with missing energy. Particularly remarkable is that this particle is a dark matter candidate, produced in roughly the right quantities in a hot early universe to account for the observed matter density. Extensive searches have been undertaken and are currently under way for such particles, through both their collision with detectors deep underground (direct detection) and their annihilations in the cosmos (indirect detection).

Another striking feature of the MSSM is the unification of the gauge couplings. For a theory with the particle content of the MSSM, assuming that all of the new particles have masses of order 1 TeV, one obtains unification of the known gauge couplings, with reasonable accuracy, at a scale $M_{\rm gut} = 2 \times 10^{16}$ GeV, corresponding to a unified coupling $\alpha_{\rm gut} \approx 1/30$. It is remarkable that these two predictions are outcomes of other requirements, and that they are consequences of symmetry.

Even before dedicated searches for supersymmetric particles were conducted at LEP, the Tevatron, and most recently the LHC, there were significant constraints on these parameters. The absence of flavor-changing neutral currents in the weak interactions of hadrons requires, in particular, a significant degree of degeneracy [or alignment (28)] in the spectrum. This degeneracy might be natural, given that in the limit of exact degeneracy, the soft parameters exhibit a significant degree of symmetry. This requires special features in the microscopic theory, which have been achieved to date only in models of gauge mediation (29) and so-called mass matrix models (30).

7. SUPERSYMMETRY: DETAILED CONSIDERATIONS OF NATURALNESS

The MSSM has provided a paradigm for experimental searches for supersymmetry, as well as theoretical efforts to construct a compelling picture of dynamical supersymmetry breaking. Notions of naturalness lead to certain expectations for the soft-breaking parameters.

For the problem of flavor, mentioned above, there are several plausible solutions. A much more serious challenge to the naturalness principle is the mass of the Higgs particle itself. Classically within the MSSM, there is a bound on the mass of the lightest Higgs particle:

$$m_H < M_Z. \tag{23}$$

This bound arises because supersymmetry strongly constrains the quartic couplings of the Higgs fields, and these are related to the gauge couplings. It turns out, however, that due to the top quark, there are significant radiative corrections to the Higgs potential (31). The diagram shown in **Figure 3**, in particular, provides a correction behaving roughly as

$$\delta\lambda = \frac{12 y_t^4}{16\pi^2} \log \frac{\tilde{m}_t}{m_t}. \tag{24}$$

More detailed studies have yielded results of the sort shown in **Figure 4**.

Figure 4 shows that, at least within the MSSM, the mass of the recently discovered Higgs particle, $m_H \approx 125$ GeV, requires that the stop be quite heavy, 8 TeV or more (alternatively, one can tune the so-called A parameter and obtain a lower stop mass). This requirement has troubling implications for naturalness. If we substitute 8 TeV on the right-hand side of Equation 21 for $\tilde{m}_t$, and take the UV cutoff to be, say, 10^{16} GeV, then the correction to the Higgs mass parameter is of order $10^4 M_Z^2$—a tuning of parameters of 1 part in 10^4.

Modifying the structure of the MSSM can help with this problem. If one adds a gauge singlet field, one can increase the quartic coupling to some degree and obtain the observed Higgs mass with significantly less [(although still appreciable (33)] tuning.

The experimental programs at LEP, the Tevatron, and the LHC have provided significant further constraints. For a broad swath of the parameter space, independent of the arguments about the Higgs mass, squark and gaugino masses are now known to be greater than 1 TeV. The

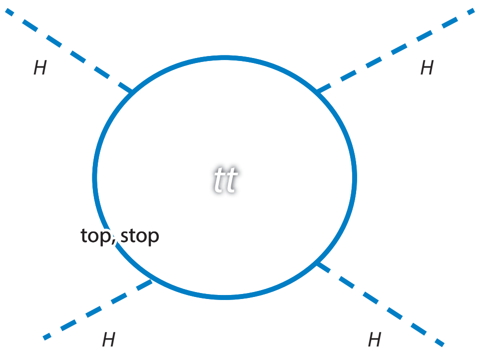

Figure 3

Top and stop corrections to the Higgs quartic coupling in the minimal supersymmetric Standard Model.

resulting contributions to the Higgs mass are large, requiring fine-tuning at greater than the 1% level, independently of what physics might account for the observed mass of the Higgs boson.

8. THE COSMOLOGICAL CONSTANT AND INFLATION

Within the framework of known physics, there is a far more serious violation of naturalness that we have not yet confronted: the size of the dark energy or cosmological constant. A cosmological constant is a dimension-zero term in the effective action that is even more problematic than the dimension-two Higgs mass term:

$$\mathcal{L}_\Lambda = \int d^4x\sqrt{g}\Lambda. \tag{25.}$$

Assuming that the observed dark energy is a cosmological constant, we have $\Lambda \approx 10^{-47}$ GeV^4. This is an extremely small number in particle physics units; absent any general principle, one might have expected $\Lambda \approx M_p^4$, roughly 120 orders of magnitude larger. As for the Higgs mass problem, this estimate is reinforced by a simple calculation. In a quantum field theory, even if the vacuum energy vanishes classically, there is a quantum contribution to the energy, which is simply

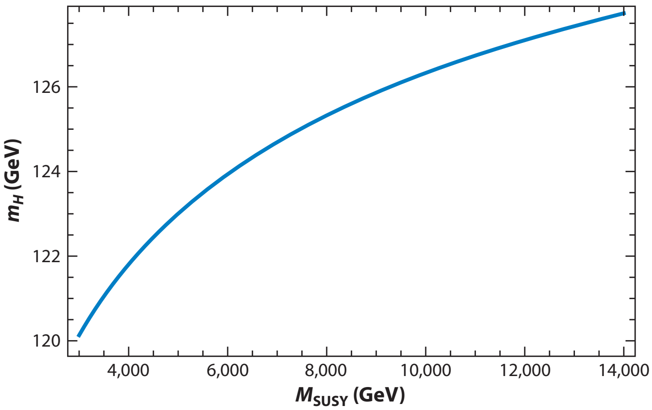

Figure 4

Higgs mass as a function of the stop mass for large tan β and a small value of the A parameter, including only leading-log corrections. More complete and detailed results can be found in, for instance, Reference 32.

a sum of the zero-point energy for bosons and the energy of the filled Dirac sea for fermions:

$$\Lambda = \sum_{\text{helicities}} (-1)^F \frac{1}{2} \int \frac{d^3k}{(2\pi)^3} \sqrt{k^2 + m^2}. \tag{26}$$

Here, $(-1)^F$ is $+1$ for bosons and -1 for fermions. Each term in the sum is quartically divergent. Taking M_p as the cutoff yields the naïve estimate.

In the case of supersymmetric theories, the situation is somewhat better. The number of bosonic and fermionic degrees of freedom is the same, and the leading divergence cancels. But one gets a result proportional to the fourth power of the supersymmetry-breaking scale. Even for the lowest conceivable supersymmetry-breaking scale (TeV), this result is many orders of magnitude larger than the observed dark energy.

In fact, there is no proposal to understand the small value of the dark energy in 't Hooftian terms; general relativity simply does not become more symmetric in the limit $\Lambda \to 0$. Calculations in string theory, the only framework we have in which dark energy may be calculable, are consistent with expectations based on dimensional analysis (34).

The value of the cosmological constant is remarkable in another way. Although small in particle physics units, it is substantial in units relevant to the present cosmological epoch; indeed, the cosmological constant has just become important "recently" (that is, the past few billion years), and it will dominate the energy density forever. One could imagine that some dynamics couples the cosmological constant and the density of dark matter, for example, but no such connection has been uncovered. Instead, Weinberg (35), following a suggestion by Banks (36) and Linde (37), proposed an explanation of a different type. He imagined that the observed Universe is part of a larger structure, subsequently dubbed a multiverse, in which the cosmological constant can take a range of values that are, essentially, randomly distributed. If one could take an inventory of this multiverse, one would find that only in some regions are there observers. This criterion, known as the anthropic principle, is much like arguing that observers (e.g., humans) are found only in a very tiny fraction of the volume of the Universe, on the surfaces of planets with liquid water.

At a minimum, Weinberg (35) argued, a universe supporting observers should contain galaxies. In our Universe, galaxies formed approximately one billion years or so after the big bang, the time required for small primordial density fluctuations (presumably formed during an epoch of inflation) to grow and become nonlinear. If the cosmological constant were so large that it dominated the energy density one billion years after the big bang, structure would not form.

An additional, crucial element of the argument relies precisely on the fact that the cosmological constant is unnatural: There is nothing more symmetric or otherwise special about a Lagrangian with vanishing Λ, so it is reasonable to expect that the probability of finding one or another value of Λ near zero is uniform. Thus, in particular, one is likely to find the largest value of Λ consistent with the anthropic requirement above. This value is somewhat larger than the dark energy that was subsequently discovered. More refined versions of the argument (38) lead to values closer to the observed value.

One may or may not be troubled by entertaining the possibility that anthropic considerations determine the laws of physics, and one can debate how significant is the success of predicting, at least at a rough order-of-magnitude level, the cosmological constant. Perhaps a more compelling concern raised by such considerations is simply whether or not there exist physical theories in which such a possibility is realized. The number of possible configurations that must be surveyed is enormous; given the small value of the cosmological constant in typical particle physics units, one might imagine that there should be at least 10^{120} such states. Several researchers have put forward scenarios in which such a landscape of possibilities, usually thought of as (metastable) vacua of some underlying theory, might arise (39–41). In string theory with some compactified dimensions, in particular, there are many types of quantized flux (analogous to magnetic flux in QED) that can take many values, providing the potential for vast numbers of possible states. In each of these vacua, the degrees of freedom and the parameters of the Lagrangian take different values. If there are enough such states, the parameters will be densely distributed. The existence of such a landscape or discretuum of vacua remains conjectural, however.

The success, to the degree that it may be counted as one, of anthropic considerations for the cosmological constant raises the possibility [or concern (?)] that such considerations might govern other features of our SMs, most notably the Higgs mass. Indeed, this mass is not nearly as severely tuned as the cosmological constant. Moreover, it is plausible that the TeV scale is anthropically selected. If the Higgs mass squared were much larger than it is, either (*a*) electroweak symmetry would be unbroken or (*b*) it would be broken and the W and Z bosons would be extremely heavy. In either case, life would likely be impossible. If stars existed at all, their properties would be quite different than those in our Universe, affecting important quantities such as the abundance of heavy elements.[3]

Thus, it is conceivable that the value of the Higgs mass is selected by anthropic considerations from a landscape of possibilities. If so, the naturalness principle might not be operative, and the value of the electroweak scale might not have any additional consequences for low-energy physics.

Other aspects of cosmology raise serious questions of naturalness as well. Inflation, the proposal that the Universe went through a period of extremely rapid expansion early in its history has received extensive experimental support in the past two decades from studies of the cosmic microwave radiation background (43). Inflation explains the homogeneity and flatness of the Universe, as well as the structure we observe about us, but existing models of the phenomenon suffer from problems of fine-tuning in varying degrees. It is plausible that anthropic considerations might play some role here as well.

9. OTHER ARENAS FOR QUESTIONS OF NATURALNESS

Another puzzling number in the SM is mentioned above: the small value of $\theta_{\rm QCD}$. Interestingly, this problem is not likely to be solved anthropically (44). Provided that θ is less than some moderately small number (0.01 or even larger), nothing changes qualitatively in the strong interactions; indeed, the dependence on θ of nuclear reaction rates, for example, is very weak (45).

Solutions that are compatible with notions of naturalness have been put forward. They rely, ultimately, on the fact that QCD, considered in isolation, becomes more symmetric in the limit $\theta \to 0$. One possibility is that the mass of the up quark is very small. In the limit $m_u \to 0$, θ is unobservable and *CP* is conserved in the strong interactions. d_n is smaller than the experimental limit, provided that

$$\frac{m_u}{m_d} < 10^{-10}. \quad 27.$$

The main question is whether a small up quark mass is compatible with properties of the strong interactions. Researchers have debated this question over the years (46–48), but lattice gauge theory calculations appear to conclusively rule out this possibility (49–51).

A second proposal involves the axion, a pseudo-Goldstone particle associated with an approximate global Peccei–Quinn symmetry. This field would couple to $F\tilde{F}$. If its potential arises only through this coupling, then it has a minimum near the origin, where the theory conserves *CP*. To ensure that the QCD corrections dominate, the Peccei–Quinn symmetry must be an extremely

[3]For a recent, wide-ranging discussion of these issues, with extensive references, see Reference 42.

good symmetry. The axion mass is of order

$$m_a^2 \approx \frac{m_\pi^2 f_\pi^2}{f_a^2}, \qquad 28.$$

which, for values of f_a of order 10^{11} GeV or larger (as required from astrophysical considerations), is extremely small. For a range of parameters (and depending on assumptions about early cosmic history), the axion can be dark matter.

As a consequence of the small axion mass, tiny, *CP*-asymmetric and Peccei–Quinn symmetry–violating couplings can give rise to an unacceptably large θ. Several proposals have been put forward to achieve a Peccei–Quinn symmetry of sufficient quality; the most compelling come from string theory (52–54). In an interesting range of its parameter space, this particle can play the role of dark matter [raising the possibility of some sort of anthropic selection for axions and, hence, small θ (55)].

A third proposal is that *CP* is conserved in the underlying theory and spontaneously broken in a way that reproduces the measured, order-one *CP*-violating phase, δ, in the Standard Model CKM matrix with a tiny θ. Models for such a phenomenon have been proposed (56, 57). However, there are many difficulties in assuring that θ remains small when higher-dimensional operators and quantum corrections are taken into account. In a landscape framework (discussed below), whereas *CP* is indeed conserved in the underlying theory, *CP*-conserving ground states (i.e., states in which the "bare" θ might be expected to be zero) are likely to be very rare. At present, then, it appears that the axion is the most plausible solution of the strong *CP* problem.

Once one has admitted the possibility of anthropic selection, one is forced to contemplate its relevance even for quantities that are naturally small. One may well imagine that anthropic considerations could play a role in determining the masses of the up and down quarks and the electron, although their possible relevance for heavier quarks and leptons is not obvious.

10. MODEL LANDSCAPES

As mentioned above, compactification of string theory with fluxes provides a model of how a landscape might arise. In interesting constructions, the number of possible flux types is often large (on the order of hundreds or more), and these fluxes can range over many discrete values. For each choice of flux, there may be many stationary points of the effective action. In this way, one can build up an exponentially large number of states, creating a setting for Weinberg's solution of the cosmological constant problem.

Establishing the existence of a discretuum poses many challenges. Before turning on fluxes, in the classical approximation, string vacua exhibit large, continuous degeneracies. Associated with these degeneracies are large numbers of scalar fields, termed moduli, without potentials. Turning on fluxes often provides potentials for many of these fields, with stable minima. But, again at the classical level, there are invariably some massless fields left over. It is plausible that some or all of these remaining fields are stabilized by nonperturbative effects. Kachru et al. (41) proposed scenarios giving rise to the existence of isolated vacua with supersymmetry or approximate supersymmetry. Actually constructing such vacua in a consistent manner is challenging; it is debatable, for example, whether there is ever a small parameter that enables systematic study.

Assuming the existence of a landscape, the interesting issue is to understand the statistics of these states. One might hope, given knowledge of the distribution of parameters and some observational or anthropic constraints, to establish that, for example, low-energy supersymmetry is or is not likely; indeed, as discussed further in Section 10, doing so would provide a quite concrete realization of notions of naturalness. Researchers have attempted to understand such statistics (58), and several plausible arguments have been put forward.

1. Among nonsupersymmetric stationary points, only a very tiny fraction is metastable. This suggests, but hardly proves, that some degree of supersymmetry might be an outcome (59).
2. Among the remaining nonsupersymmetric states, with a small cosmological constant, the vast majority are short-lived (60, 61).
3. Among supersymmetric states, if supersymmetry is not dynamically broken, high scales of supersymmetry breaking are favored (62, 63). With dynamical breaking, lower scales may be favored.
4. As discussed further in Section 12, below, states exhibiting certain types of (ordinary) symmetries are rare.

Even in the absence of a completely reliable model, if we assume the existence of a landscape, many of these features seem robust. They rely on quite minimal assumptions about the features of low-energy effective actions and distributions of Lagrangian parameters.

11. NATURALNESS IN A LANDSCAPE FRAMEWORK

Above, I present the rather bleak prospect that certain parameters of the SM, such as the Higgs mass, are completely determined by anthropic considerations, and that considerations of naturalness—and with them interesting possibilities for new TeV-scale degrees of freedom and new symmetries—play no role. But there are intermediate possibilities, which should be considered with greater care.

Indeed, a landscape in some sense provides an ideal setting in which to consider questions of naturalness and to understand how it might emerge, sometimes or always, as a governing principle. Weinberg's cosmological constant argument relies crucially on the assumption that there is nothing special, at a fundamental level, about the point where $\Lambda = 0$. For the Higgs boson, things might be different if nature is approximately supersymmetric. Indeed, in studies of model landscapes (58, 62, 63), several branches have been identified that differ in the nature of the realization of supersymmetry.

1. A nonsupersymmetric branch, where the distribution of the Higgs mass squared is roughly uniform. The cost of having a Higgs mass m_H is m_H^2/M_p^2.
2. A supersymmetric branch where the breaking of supersymmetry is nondynamical. Here, supersymmetry, despite the cancellation of quadratic divergences, does not help; the fraction of states with larger breaking of supersymmetry, F, grows as a large power of F. So, operationally, this branch is like branch 1.
3. A supersymmetric branch with dynamically broken supersymmetry (in the sense that the supersymmetry-breaking scale behaves as $e^{-\frac{8\pi^2}{bg^2}}$). Here, the number of states with a small Higgs mass and a small cosmological constant is the same per decade as a function of the supersymmetry-breaking scale. Without the introduction of other considerations (perhaps the density of dark matter), there is no preference for low-scale supersymmetry breaking. Conceivably, other such considerations would favor a scale more like 8 TeV than 1 TeV.
4. A supersymmetric branch favoring low-scale supersymmetry breaking, in which other quantities (the value of the superpotential and the μ parameter) are dynamically determined as well (corresponding to dynamical breaking of so-called R symmetries). Here, the lowest possible scale of supersymmetry breaking is favored. General arguments can be put forward suggesting that there are far fewer states on this branch than on branch 3.

12. HOW NATURAL ARE SYMMETRIES?

In a landscape framework, one can revisit the question of symmetries themselves. The symmetries one has in mind are discrete symmetries, global continuous symmetries, gauge symmetries, and supersymmetry. 't Hooft's naturalness principle assumes that symmetries themselves are special or singled out. Of these various types of symmetries, global continuous symmetries are not a feature of quantum gravity theories [in string theories, this statement is often a theorem (64)]. Gauge symmetries appear common in string theories, as does supersymmetry. Discrete symmetries appear frequently as well. It is these latter symmetries that are of particular interest. They might account for the stability of the proton in supersymmetric theories and the smallness of the Yukawa couplings of the SM, and in the construction of particle physics models it is usually assumed that they are somehow singled out. Yet, in the flux landscapes that have been studied, states (vacua) with symmetries would appear to be quite rare (65).

To understand this apparent rarity, one can ask how symmetries arise when one compactifies a theory on some compact space. In many solutions of string theory, the compact space exhibits discrete symmetries. These are typically subgroups of the original rotational symmetry of the higher-dimensional space. These symmetries translate into conventional discrete symmetries of the field theory that describes the system at low energies. Now imagine turning on fluxes. Typical fluxes will transform under these rotations; as a result, the low-energy theory does not exhibit the symmetry. Recall that in the flux landscape, the large number of states results from the large number of possible fluxes. If most of the fluxes are not invariant under the symmetry—the typical situation—then at best an exponentially small fraction of the states will exhibit the symmetry. These considerations apply to the sorts of discrete symmetries we might invoke to explain proton stability, as well as to *CP*. There may be other (cosmological?) considerations that would favor symmetric states (66), but this simple observation calls into question the assumptions underlying 't Hooft's naturalness criterion.

Interestingly, supersymmetry might function differently. Another issue related to landscapes is stability; a state with a small cosmological constant, similar to our own, will be surrounded by vast numbers of states with negative cosmological constant. It is necessary that the lifetime for decay of the state to every single one of its neighbors be extremely small (60, 61). It turns out that the simplest way to account for such stability involves approximate supersymmetry of the state. In the limit of exact supersymmetry, in fact, the symmetry insures exact stability; if the breaking is small, the lifetime of the state becomes exponentially long as the breaking scale becomes small (60).

13. CONCLUSIONS: NATURALNESS AS A GUIDE

It is still possible that nature is "natural" in the sense of 't Hooft. Future runs of the LHC might provide evidence for supersymmetry, warped extra dimensions, or some variant of technicolor. But the current experimental situation raises the unsettling possibility that naturalness may not be a good guiding principle. Indeed, naturalness is in tension with another principle: simplicity. Simplicity has a technical meaning: The simplest theory is the one with the fewest degrees of freedom consistent with known facts. Contrast, for example, the minimal SM, including its single Higgs doublet, with supersymmetric theories, including their many additional fields and couplings. So far, the experimental evidence suggests that simplicity is winning. The observed Higgs mass is in tension with expectations from supersymmetric theories, but also with technicolor and other proposals.

The main alternative to natural theories (apart from the possibility that extreme fine-tuning is simply a fact) is the landscape or multiverse. In such a situation, our neighborhood in the

Universe might be simple, but the underlying structure is unimaginably complex. As discussed above, however, this idea has at least one major success: the prediction of dark energy. This prediction provides a plausible picture for other (but not necessarily all) tuned quantities.

Why might we subscribe to a naturalness principle? After all, if the Universe is described by a single theory, with a single set of degrees of freedom and a single Lagrangian with fixed parameters, the question of fine-tuning is metaphysical; things are the way they are, and it is not clear why we should be troubled with the value of some parameter or other. The landscape has the potential to make the question of naturalness concrete. A theory (a set of degrees of freedom and Lagrangian parameters) is natural in a landscape if it is typical of the states compatible with features observed in nature (e.g., small cosmological constant, large hierarchy). We have seen that model landscapes may prefer, for example, no supersymmetry or a very high scale of supersymmetry breaking. Conventional symmetries, such as discrete symmetries (including *CP*), would seem rare. Alternatively, this review presents arguments that in a landscape supersymmetric states might be common, and that classes of these states would favor supersymmetry in the conventional way. It is possible that the next round of LHC experiments or experiments at slightly higher energies will discover evidence for supersymmetry, large extra dimensions, or totally unanticipated phenomena that will restore our confidence in the notion of naturalness.

DISCLOSURE STATEMENT

The author is not aware of any affiliations, memberships, funding, or financial holdings that might be perceived as affecting the objectivity of this review.

ACKNOWLEDGMENTS

The writing of this review was supported by the US Department of Energy, grant number DE-FG02-04ER41286. I thank many colleagues for discussions of these issues, particularly Tom Banks, Savas Dimopoulos, Patrick Draper, and Nima Arkani-Hamed.

LITERATURE CITED

1. Planck M. *Sitz. K. Preuß. Akad. Wiss. Berlin* 18:5:440 (1899)
2. Lorentz HA. *Arch. Néerl. Sci. Exactes Nat.* 25:363 (1892)
3. Weisskopf VF. *Z. Phys.* 89:27 (1934)
4. Weisskopf VF. *Z. Phys.* 90:53 (1934)
5. 't Hooft G. *NATO Sci. Ser. B* 59:135 (1980)
6. Wilson KG. *Phys. Rev. D* 3:1818 (1971)
7. Witten E. *Phys. Lett. B* 86:283 (1979)
8. Crewther RJ, Di Vecchia P, Veneziano G, Witten E. *Phys. Lett. B* 88:123 (1979)
9. Weinberg S. *Phys. Rev. D* 13:974 (1976)
10. Susskind L. *Phys. Rev. D* 20:2619 (1979)
11. Peskin ME, Takeuchi T. *Phys. Rev. Lett.* 65:964 (1990)
12. Kaplan DB, Georgi H, Dimopoulos S. *Phys. Lett. B* 136:187 (1984)
13. Georgi H, Kaplan DB. *Phys. Lett. B* 145:216 (1984)
14. Arkani-Hamed N, Cohen AG, Georgi H. *Phys. Lett. B* 513:232 (2001)
15. Arkani-Hamed N, Dimopoulos S, Dvali GR. *Phys. Lett. B* 429:263 (1998)
16. Antoniadis I, Arkani-Hamed N, Dimopoulos S, Dvali GR. *Phys. Lett. B* 436:257 (1998)
17. Arkani-Hamed N, Dimopoulos S, Dvali GR. *Phys. Rev. D* 59:086004 (1999)
18. Hoyle CD, et al. *Phys. Rev. D* 70:042004 (2004)
19. Adelberger EG, et al. *Phys. Rev. Lett.* 98:131104 (2007)

20. Randall L, Sundrum R. *Phys. Rev. Lett.* 83:3370 (1999)
21. Randall L, Sundrum R. *Phys. Rev. Lett.* 83:4690 (1999)
22. Witten E. *Nucl. Phys. B* 188:513 (1981)
23. Grisaru MT, Siegel W, Rocek M. *Nucl. Phys. B* 159:429 (1979)
24. Freedman DZ, Van Proeyen A. *Supergravity*. Cambridge, UK: Cambridge Univ. Press (2012)
25. Dine M. *Supersymmetry and String Theory: Beyond the Standard Model*. Cambridge, UK: Cambridge Univ. Press (2007)
26. Girardello L, Grisaru MT. *Nucl. Phys. B* 194:65 (1982)
27. Dimopoulos S, Georgi H. *Nucl. Phys. B* 193:150 (1981)
28. Nir Y, Seiberg N. *Phys. Lett. B* 309:337 (1993)
29. Giudice GF, Rattazzi R. *Phys. Rep.* 322:419 (1999)
30. Leurer M, Nir Y, Seiberg N. *Nucl. Phys. B* 398:319 (1993)
31. Haber HE, Hempfling R. *Phys. Rev. Lett.* 66:1815 (1991)
32. Draper P, Lee G, Wagner CEM. *Phys. Rev. D* 89:055023 (2014)
33. Hall LJ, Pinner D, Ruderman JT. *J. High Energy Phys.* 1204:131 (2012)
34. Rohm R. *Nucl. Phys. B* 237:553 (1984)
35. Weinberg S. *Phys. Rev. Lett.* 59:2607 (1987)
36. Banks T. *Nucl. Phys. B* 249:332 (1985)
37. Linde AD. *Phys. Lett. B* 175:395 (1986)
38. Weinberg S. *Phys. Rev. D* 61:103505 (2000)
39. Banks T, Dine M, Seiberg N. *Phys. Lett. B* 273:105 (1991)
40. Bousso R, Polchinski J. *J. High Energy Phys.* 0006:006 (2000)
41. Kachru S, Kallosh R, Linde AD, Trivedi SP. *Phys. Rev. D* 68:046005 (2003)
42. Hall LJ, Pinner D, Ruderman JT. arXiv:1409.0551 [hep-ph] (2014)
43. Ade PAR, et al. (Planck Collab.) arXiv:1303.5082 [astro-ph] (2013)
44. Banks T, Dine M, Gorbatov E. *J. High Energy Phys.* 0408:058 (2004)
45. Ubaldi L. *Phys. Rev. D* 81:025011 (2010)
46. Georgi H, McArthur IN. *Instantons and the up quark mass*. HUTP-81/A011. Unpubl. ms. (1981)
47. Kaplan DB, Manohar AV. *Phys. Rev. Lett.* 56:2004 (1986)
48. Banks T, Nir Y, Seiberg N. arXiv:hep-ph/9403203 (1994)
49. Brambilla N, et al. arXiv:1404.3723 [hep-ph] (2014)
50. Aoki S, et al. arXiv:1310.8555 [hep-lat] (2013)
51. Durr S, et al. *Phys. Lett. B* 701:265 (2011)
52. Witten E. *Phys. Lett. B* 149:351 (1984)
53. Bobkov K, Braun V, Kumar P, Raby S. *J. High Energy Phys.* 1012:056 (2010)
54. Dine M, Festuccia G, Kehayias J, Wu W. *J. High Energy Phys*. 1101:012 (2011)
55. Carpenter LM, Dine M, Festuccia G, Ubaldi L. *Phys. Rev. D* 80:125023 (2009)
56. Nelson AE. *Phys. Lett. B* 136:387 (1984)
57. Barr SM. *Phys. Rev. Lett.* 53:329 (1984)
58. Denef F, Douglas MR. *J. High Energy Phys.* 0405:072 (2004)
59. Marsh D, McAllister L, Wrase T. *J. High Energy Phys.* 1203:102 (2012)
60. Dine M, Festuccia G, Morisse A. *J. High Energy Phys.* 0909:013 (2009)
61. Greene B, et al. *Phys. Rev. D* 88:026005 (2013)
62. Denef F, Douglas MR. *J. High Energy Phys.* 0503:061 (2005)
63. Dine M, O'Neil D, Sun Z. *J. High Energy Phys.* 0507:014 (2005)
64. Banks T, Dixon LJ. *Nucl. Phys. B* 307:93 (1988)
65. Dine M, Sun Z. *J. High Energy Phys.* 0601:129 (2006)
66. Dine M, Festuccia G, Morisse A. *J. High Energy Phys.* 0912:013 (2009)