Unanswered Questions in the Electroweak Theory

Chris Quigg

Theoretical Physics Department, Fermi National Accelerator Laboratory, Batavia, Illinois 60510
Institut für Theoretische Teilchenphysik, Universität Karlsruhe, D-76128 Karlsruhe, Germany
Theory Group, Physics Department, CERN, CH-1211 Geneva 23, Switzerland

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Abstract

This article is devoted to the status of the electroweak theory on the eve of experimentation at CERN's Large Hadron Collider (LHC). A compact summary of the logic and structure of the electroweak theory precedes an examination of what experimental tests have established so far. The outstanding unconfirmed prediction is the existence of the Higgs boson, a weakly interacting spin-zero agent of electroweak symmetry breaking and the giver of mass to the weak gauge bosons, the quarks, and the leptons. General arguments imply that the Higgs boson or other new physics is required on the 1-TeV energy scale.

Even if a "standard" Higgs boson is found, new physics will be implicated by many questions about the physical world that the Standard Model cannot answer. Some puzzles and possible resolutions are recalled. The LHC moves experiments squarely into the 1-TeV scale, where answers to important outstanding questions will be found.

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1. INTRODUCTION

The electroweak theory (1–3) joins electromagnetism with the weak force in a single relativistic quantum field theory. Electromagnetism is a force of infinite range, whereas the influence of the charged-current weak interaction responsible for radioactive beta decay only spans distances shorter than approximately 10^{-15} cm, less than 1% of the proton radius. The two interactions, so different in their range and apparent strength, are ascribed to a common gauge symmetry. We say that the electroweak gauge symmetry is spontaneously broken to the gauge symmetry of electromagnetism. Throughout the past two decades, precision experiments have elevated the electroweak theory from a promising description to a provisional law of nature, tested as a quantum field theory at the level of 1% or better by many measurements. Joined with quantum chromodynamics (QCD)—the theory of the strong interactions—to form the Standard Model, and augmented to incorporate neutrino masses and lepton mixing, the electroweak theory describes a vast array of experimental information. The development and validation of the Standard Model are a landmark in the history of science.

Charged current: the weak interaction mediated by the $W^{\pm}$ boson, first observed in nuclear beta decay

One measure of the sweep of the electroweak theory is that its predictions hold over a prodigious range of distances, from approximately 10^{-18} m to more than 10^{8} m. The origins of the theory lie in the discovery of Coulomb's law in tabletop experiments by Cavendish (4) and Coulomb (5). It was stretched to longer and shorter distances by the progress of experiment. In the long-distance limit, the classical electrodynamics of a massless photon suffices. At shorter distances than the human scale, classical electrodynamics was superseded by quantum electrodynamics (QED), which is now subsumed in the electroweak theory that has been tested at energies up to a few hundred gigaelectronvolts. Some key steps in the evolution may be traced in References 6 and 7.

The electroweak theory anticipated the existence and properties of weak neutral-current interactions, predicted the properties of the gauge bosons $W^{\pm}$ and Z^0 that mediate charged-current and neutral-current interactions, and required the fourth quark flavor, charm. Fits to a universe of electroweak precision measurements are in excellent agreement with the Standard Model.

How the electroweak gauge symmetry is spontaneously broken is one of the most urgent and challenging questions in particle physics. The Standard-Model answer is an elementary scalar field whose self-interactions select a vacuum state in which the full electroweak symmetry is hidden. However, the elementary scalar field, also known as the Higgs boson, has not been observed directly, and we do not know whether a fundamental Higgs field exists or whether a different agent breaks electroweak symmetry. Finding the Higgs boson or its replacement is one of the great campaigns now under way in both experimental and theoretical particle physics.

The aim of this article is to survey what we know and what we need to know about the electroweak theory in anticipation of the experiments soon to begin at the Large Hadron Collider (LHC), a high-luminosity proton-proton machine that will reach 14 TeV center-of-mass (c.m.) energy. I begin in Section 2 with a short summary of the essential elements of the electroweak theory. Next, in Section 3, I examine the experimental support that has helped to establish the electroweak theory. The evidence includes the behavior of the couplings at the Lagrangian level, along with signs for weak electromagnetic unification. A prominent feature of the electroweak theory is the absence of flavor-changing neutral currents (FCNC). An important chapter in the weak interactions, just concluded, validated the picture of three-family quark mixing that organizes a vast amount of experimental information, including the observations of *CP* violation. Quantum corrections test the electroweak theory as a quantum field theory and give evidence for the interactions of (something resembling) the Higgs boson with the weak gauge bosons. Low-energy tests of the electroweak theory can be expressed as determinations of the weak mixing parameter. The electroweak theory gives but a partial explanation for the origin of quark and lepton masses, so I regard all the quark and lepton masses as evidence for physics beyond the Standard Model.

The Higgs boson, the missing ingredient of the Standard Model, is the subject of Section 4. There I describe theoretical and experimental constraints on the Higgs boson mass and outline the production and decay characteristics that will govern the search at the LHC. Alternatives to the Higgs mechanism, beginning with dynamical symmetry breaking inspired by the microscopic theory of the superconducting phase transition, are described. I devote a brief passage to what the world would have been like in the absence of an explicit mechanism to hide the electroweak symmetry. This excursion underscores the importance of discovering the agent of electroweak symmetry breaking for our understanding of the everyday world.

Section 5 is devoted to the shortcomings of the Standard Model, which include the partial understanding of fermion masses and mixing among quark families, the challenge of stabilizing the Higgs mass below 1 TeV in the face of quantum corrections, and the vacuum energy problem. I take note of questions that lie beyond the scope of the Standard Model: the nature of dark matter, the matter asymmetry of the universe, the quantization of electric charge, and the role of gravity. Both sets of issues motivate more complete and predictive extensions to the Standard Model.

The new era ushered in by the LHC is the subject of Section 6. I pose a series of electroweak questions for the LHC, and then note some possibilities for new physics motivated by the hierarchy

Neutral current: the weak interaction mediated by the Z^0 boson, first observed in the reactions $\nu_\mu e \to \nu_\mu e$ and $\nu_\mu N \to \nu_\mu +$ anything

Higgs boson: elementary scalar particle that is the avatar of electroweak symmetry breaking in the Standard Model; awaits discovery

Large Hadron Collider (LHC): research instrument at CERN; 27 km in circumference. It is intended to provide proton-proton collisions at 14-TeV c.m. energy and luminosity exceeding 10^{34} cm^{-2} s^{-1}, as well as collisions between heavy ions. See **http://lhc.web.cern.ch**

Flavor-changing neutral current (FCNC): a transition that changes quark or lepton flavor, without changing electric charge; strongly inhibited by the Glashow, Iliopoulos & Maiani (53) mechanism in the standard electroweak theory

Dynamical symmetry breaking: occurs when the spontaneous breakdown of a symmetry is a consequence of strong forces among the constituents, leading to composite Goldstone bosons problem and the search for dark matter candidates. I describe how new knowledge might build up as the LHC data samples grow, and remark on the continuing role of experiments at the intensity frontier. A short summary concludes the article in Section 7.

2. THE ELECTROWEAK THEORY

The electroweak theory, and the path by which it evolved, is developed in many modern textbooks, some of which are cited at the end of this article (see Related Resources). Useful perspectives on the current situation are presented in lecture courses, including References 8 through 10. Here I give a brief summary of the essential ideas and outcomes.

We build the Standard Model of particle physics on a set of constituents that we regard provisionally as elementary: the quarks and leptons, plus a few fundamental forces derived from gauge symmetries. The quarks are influenced by the strong interaction and so carry color, the strong-interaction charge, whereas the leptons do not feel the strong interaction and are colorless. We idealize the quarks and leptons as pointlike because they show no evidence of internal structure at the current limit of our resolution, $r \lesssim 10^{-18}$ m. The charged-current weak interaction responsible for radioactive beta decay and other processes acts only on the left-handed fermions. Whether the observed parity violation reflects a fundamental asymmetry in the laws of nature or a left-right symmetry that is hidden by circumstance and might be restored at higher energies is uncertain.

Like its forerunner, QED, the electroweak theory is a gauge theory, in which interactions follow from symmetries. The correct electroweak gauge symmetry, which melds an $SU(2)_{\mathrm{L}}$ family (weak-isospin) symmetry with a $U(1)_Y$ weak-hypercharge phase symmetry, emerged through trial and error, guided by experiment. We characterize the leptonic sector of the $SU(2)_{\mathrm{L}} \otimes U(1)_Y$ theory by the left-handed leptons

$$\mathsf{L}_e = \begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_{\mathrm{L}}, \quad \mathsf{L}_\mu = \begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_{\mathrm{L}}, \quad \text{and} \quad \mathsf{L}_\tau = \begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}_{\mathrm{L}}, \qquad 1.$$

with weak isospin $I = \frac{1}{2}$ and weak hypercharge $Y(\mathsf{L}_\ell) = -1$, and the right-handed weak-isoscalar charged leptons

$$\mathsf{R}_{e,\mu,\tau} = e_{\mathrm{R}}, \mu_{\mathrm{R}}, \tau_{\mathrm{R}}, \qquad 2.$$

with weak hypercharge $Y(\mathsf{R}_\ell) = -2$. The weak hypercharges are chosen to reproduce the observed electric charges through the connection $Q = I_3 + \frac{1}{2}Y$. Here we have idealized the neutrinos as massless. Very brief comments on massive neutrinos can be found in Section 5.1.

The hadronic sector consists of the left-handed quarks

$$\mathsf{L}_q^{(1)} = \begin{pmatrix} u \\ d' \end{pmatrix}_{\mathrm{L}}, \quad \mathsf{L}_q^{(2)} = \begin{pmatrix} c \\ s' \end{pmatrix}_{\mathrm{L}}, \quad \text{and} \quad \mathsf{L}_q^{(3)} = \begin{pmatrix} t \\ b' \end{pmatrix}_{\mathrm{L}}, \qquad 3.$$

with weak isospin $I = \frac{1}{2}$ and weak hypercharge $Y(\mathsf{L}_q) = \frac{1}{3}$, and their right-handed weak-isoscalar counterparts

$$\mathsf{R}_u^{(1,2,3)} = u_{\mathrm{R}}, c_{\mathrm{R}}, t_{\mathrm{R}} \quad \text{and} \quad \mathsf{R}_d^{(1,2,3)} = d_{\mathrm{R}}, s_{\mathrm{R}}, b_{\mathrm{R}}, \qquad 4.$$

with weak hypercharges $Y(\mathsf{R}_u) = \frac{4}{3}$ and $Y(\mathsf{R}_d) = -\frac{2}{3}$. The primes on the lower components of the quark doublets in Equation 3 signal that the weak eigenstates are mixtures of the mass eigenstates:

$$\begin{pmatrix} d' \\ s' \\ b' \end{pmatrix} = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} \begin{pmatrix} d \\ s \\ b \end{pmatrix} \equiv \mathsf{V} \begin{pmatrix} d \\ s \\ b \end{pmatrix}, \qquad 5.$$

where the 3×3 unitary Cabibbo (11)–Kobayashi–Maskawa (12) matrix V, colloquially termed the CKM matrix, expresses the quark mixing. See Section 3.3 for further discussion. The fact that each left-handed lepton doublet is matched by a left-handed quark doublet guarantees that the theory is anomaly free, so that quantum corrections respect the gauge symmetry (13).

The $SU(2)_L \otimes U(1)_Y$ electroweak gauge group entails two sets of gauge fields: (*a*) a weak isovector $\boldsymbol{b}_\mu$, with coupling constant g, and (*b*) a weak isoscalar $\mathcal{A}_\mu$, with its own coupling constant g'. The gauge fields compensate for the variations induced by gauge transformations, provided that they obey the transformation laws $\boldsymbol{b}_\mu \to \boldsymbol{b}_\mu - \boldsymbol{\alpha} \times \boldsymbol{b}_u - (1/g)\partial_\mu \boldsymbol{\alpha}$ under an infinitesimal weak-isospin rotation generated by $G = 1 + (i/2)\boldsymbol{\alpha} \cdot \boldsymbol{\tau}$ (where $\boldsymbol{\tau}$ are the Pauli isospin matrices) and $\mathcal{A}_\mu \to \mathcal{A}_\mu - (1/g')\partial_\mu \alpha$ under an infinitesimal hypercharge phase rotation. Corresponding to these gauge fields are the field-strength tensors

$$F^\ell_{\mu\nu} = \partial_\nu b^\ell_\mu - \partial_\mu b^\ell_\nu + g\varepsilon_{jk\ell} b^j_\mu b^k_\nu, \qquad 6.$$

where $\ell = 1, 2, 3$ for the weak-isospin symmetry, and

$$f_{\mu\nu} = \partial_\nu \mathcal{A}_\mu - \partial_\mu \mathcal{A}_\nu \qquad 7.$$

for the weak-hypercharge symmetry.

We may summarize the interactions by the Lagrangian

$$\mathcal{L} = \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{leptons}} + \mathcal{L}_{\text{quarks}}, \qquad 8.$$

with

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4}\sum_\ell F^\ell_{\mu\nu} F^{\ell\mu\nu} - \frac{1}{4} f_{\mu\nu} f^{\mu\nu}, \qquad 9.$$

$$\begin{aligned}\mathcal{L}_{\text{leptons}} = {}& \bar{\mathsf{R}}_\ell i\gamma^\mu \left(\partial_\mu + i\frac{g'}{2}\mathcal{A}_\mu Y\right)\mathsf{R}_\ell \\ &+ \bar{\mathsf{L}}_\ell i\gamma^\mu \left(\partial_\mu + i\frac{g'}{2}\mathcal{A}_\mu Y + i\frac{g}{2}\boldsymbol{\tau}\cdot\boldsymbol{b}_\mu\right)\mathsf{L}_\ell,\end{aligned} \qquad 10.$$

where ℓ runs over e, μ, τ, and

$$\begin{aligned}\mathcal{L}_{\text{quarks}} = {}& \bar{\mathsf{R}}^{(n)}_u i\gamma^\mu \left(\partial_\mu + i\frac{g'}{2}\mathcal{A}_\mu Y\right)\mathsf{R}^{(n)}_u \\ &+ \bar{\mathsf{R}}^{(n)}_d i\gamma^\mu \left(\partial_\mu + i\frac{g'}{2}\mathcal{A}_\mu Y\right)\mathsf{R}^{(n)}_d \\ &+ \bar{\mathsf{L}}^{(n)}_q i\gamma^\mu \left(\partial_\mu + i\frac{g'}{2}\mathcal{A}_\mu Y + i\frac{g}{2}\boldsymbol{\tau}\cdot\boldsymbol{b}_\mu\right)\mathsf{L}^{(n)}_q,\end{aligned} \qquad 11.$$

where the generation index n runs over 1, 2, 3.

Although the weak and electromagnetic interactions share a common origin in the $SU(2)_L \otimes U(1)_Y$ gauge symmetry, their manifestations are very different. Electromagnetism is a force of infinite range, whereas the influence of the charged-current weak interaction responsible for radioactive beta decay only spans distances shorter than approximately 10^{-15} cm. The established phenomenology of the weak interactions is thus at odds with the theory we have developed to this point. The gauge Lagrangian (Equation 9) contains four massless electroweak gauge bosons, namely $b^1_\mu, b^2_\mu, b^3_\mu$, and $\mathcal{A}_\mu$. They are massless because a mass term such as $\frac{1}{2}m^2\mathcal{A}_\mu\mathcal{A}^\mu$ is not invariant under a gauge transformation. Nature has but one: the photon. Moreover, the $SU(2)_L \otimes U(1)_Y$ gauge symmetry forbids fermion mass terms $m\bar{f}f = m(\bar{f}_R f_L + \bar{f}_L f_R)$ in Equations 10 and 11 because the left-handed and right-handed fields transform differently.

Cabibbo–Maskawa–Kobayashi (CKM) matrix: the quark-mixing matrix (Equation 5) that maps mass eigenstates to flavor eigenstates, so called to recognize both Nicola Cabibbo's formulation of hadron-lepton universality for charged-current interactions (11) and the observation by Makoto Kobayashi and Toshihide Maskawa that mixing among three quark doublets yields a nontrivial phase that could account for *CP* violation (12)

Anomaly: the violation by quantum corrections of a symmetry of the Lagrangian. If anomalies violate gauge symmetry, the theory becomes inconsistent, so the freedom from anomalies becomes a powerful condition on candidate theories, as emphasized for the electroweak theory (13)

To give masses to the gauge bosons and constituent fermions, we must hide the electroweak symmetry, recognizing that a symmetry of the laws of nature does not imply that the same symmetry is manifest in the outcomes of those laws. The superconducting phase transition offers an instructive model for hiding the electroweak gauge symmetry. To give masses to the intermediate bosons of the weak interaction, we appeal to the Meissner effect—the exclusion of magnetic fields from a superconductor, which corresponds to the photon developing a nonzero mass within the superconducting medium. What has come to be referred to as the Higgs mechanism (14–17) can be understood as a relativistic generalization of the Ginzburg–Landau phenomenology (18) of superconductivity.

Let us see how spontaneous symmetry breaking operates in the electroweak theory. We introduce a complex doublet of scalar fields

$$\phi \equiv \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} \qquad 12.$$

with weak hypercharge $Y_\phi = +1$. Next, we add to the Lagrangian new (gauge-invariant) terms for the interaction and propagation of the scalars,

$$\mathcal{L}_{\text{scalar}} = (\mathcal{D}^\mu \phi)^\dagger (\mathcal{D}_\mu \phi) - V(\phi^\dagger \phi), \qquad 13.$$

where the gauge-covariant derivative is

$$\mathcal{D}_\mu = \partial_\mu + i\frac{g'}{2}\mathcal{A}_\mu Y + i\frac{g}{2}\tau \cdot \boldsymbol{b}_\mu \qquad 14.$$

and [inspired by Ginzburg & Landau (18)] the potential interaction has the form

$$V(\phi^\dagger \phi) = \mu^2(\phi^\dagger \phi) + |\lambda|(\phi^\dagger \phi)^2. \qquad 15.$$

We are also free to add gauge-invariant Yukawa interactions between the scalar fields and the leptons (ℓ runs over e, μ, τ as before),

$$\mathcal{L}_{\text{Yukawa}-\ell} = -\zeta_\ell[(\bar{\mathsf{L}}_\ell \phi)\mathsf{R}_\ell + \bar{\mathsf{R}}_\ell(\phi^\dagger \mathsf{L}_\ell)], \qquad 16.$$

and similar interactions with the quarks.

Now we arrange the self-interactions of the scalars so that the vacuum state corresponds to a broken-symmetry solution. The electroweak symmetry is spontaneously broken if the parameter μ^2 is taken to be negative. In that event, gauge invariance gives us the freedom to choose the state of minimum energy—the vacuum state—to correspond to the vacuum expectation value

$$\langle \phi \rangle_0 = \begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix}, \qquad 17.$$

where $v = \sqrt{-\mu^2/|\lambda|}$.

The vacuum of Equation 17 breaks the gauge symmetry $SU(2)_{\text{L}} \otimes U(1)_Y \rightarrow U(1)_{\text{em}}$. The vacuum state $\langle \phi \rangle_0$ is invariant under a symmetry operation corresponding to the generator $\mathcal{G}$ provided that $e^{i\alpha\mathcal{G}} \langle \phi \rangle_0 = \langle \phi \rangle_0$, that is, if $\mathcal{G} \langle \phi \rangle_0 = 0$. Direct calculation reveals that the original four generators are all broken but that electric charge is not. The photon remains massless, but the other three gauge bosons acquire masses, as auxiliary scalars assume the role of the third (longitudinal) degrees of freedom.

By introducing the weak mixing angle θ_{W} through the definition $g' = g \tan \theta_{\text{W}}$, we can express the photon as the linear combination $A = \mathcal{A} \cos \theta_{\text{W}} + b_3 \sin \theta_{\text{W}}$. We identify the strength of its (pure vector) coupling to charged particles, $gg'/\sqrt{g^2 + g'^2}$, with the electric charge e. The mediator of the charged-current weak interaction, $W^{\pm} = (b_1 \mp i b_2)/\sqrt{2}$, acquires a mass

$$M_W = gv/2 = ev/2\sin\theta_W. \tag{18.}$$

The electroweak gauge theory reproduces the low-energy phenomenology of the $V - A$ theory of weak interactions, provided that we set $v = (G_F\sqrt{2})^{-1/2} = 246$ GeV, where $G_F = 1.16637(1) \times 10^{-5}$ GeV^{-2} is Fermi's weak-interaction coupling constant. It follows at once that $M_W \approx 37.3$ GeV$/\sin\theta_W$. The combination of the I_3 and Y gauge bosons orthogonal to the photon is the mediator of the neutral-current weak interaction, $Z = b_3\cos\theta_W - \mathcal{A}\sin\theta_W$, which acquires a mass

$$M_Z = M_W/\cos\theta_W. \tag{19.}$$

The masses of the elementary fermions are not predicted by the electroweak theory. Each fermion mass involves a new Yukawa coupling ζ (cf. Equation 16). When the electroweak symmetry is spontaneously broken, the electron mass emerges as $m_e = \zeta_e v/\sqrt{2}$. The Yukawa couplings that reproduce the observed quark and lepton masses range over many orders of magnitude, as detailed in Section 3.7. We do not know what sets the values of the Yukawa couplings. They do not follow from, for example, a known symmetry principle.

Three of the four scalar degrees of freedom that we introduced to contrive a vacuum state that hides the electroweak gauge symmetry have become the longitudinal components of W^+, W^-, and Z. The fourth appears as a massive spin-zero particle known as the Higgs boson, H, a vestige of the spontaneous symmetry breaking. Its mass is given symbolically as $M_H^2 = -2\mu^2 > 0$, but we have no prediction for its value. However, the interactions of the Higgs boson with gauge bosons and fermions are completely specified—after spontaneous symmetry breaking—by the Lagrangian terms $\mathcal{L}_{\text{scalar}}$ and $\mathcal{L}_{\text{Yukawa}}$. Given the mass of the Higgs boson, we may calculate its properties.

Let us summarize how particle mass arises in the standard electroweak theory. Unless the electroweak gauge symmetry is hidden, the four gauge bosons and all the constituent fermions are massless. Spontaneous symmetry breaking, in the form of the Higgs mechanism, gives masses to the weak gauge bosons and creates the possibility for the fermions to acquire mass. Once the weak mixing parameter $\sin^2\theta_W$ is fixed by the study of weak neutral-current interactions, the theory makes successful quantitative predictions for the $W^{\pm}$ and Z boson masses. Although the natural scale of fermion masses would seem to be set by the electroweak scale, the specific values are determined by Yukawa couplings of the fermions to the Higgs field. These Yukawa couplings are not predicted by the electroweak theory. Finally, the theory requires a scalar Higgs boson but does not make an explicit prediction for its mass.

3. HOW THE ELECTROWEAK THEORY BECAME A LAW OF NATURE… AND WHAT WE REALLY KNOW

The $SU(2)_L \otimes U(1)_Y$ electroweak theory was formulated in the context of extensive experimental information about the charged-current weak interactions. Central elements included the parity-violating $V - A$ structure of the charged current and the Cabibbo universality of leptonic and semileptonic processes. On the theoretical front, a classic unitarity argument (19) made it clear that Fermi's four-fermion description could not be valid above c.m. energy $\sqrt{s} = 620$ GeV. Analysis of the reaction $\nu\bar{\nu} \to W^+W^-$ showed that the ad hoc introduction of intermediate vector bosons, to make the weak interaction nonlocal, had divergence diseases of its own (20).

The weak neutral-current interaction was not detected before the electroweak theory was formulated. The prediction of this new phenomenon and the availability of high-energy neutrino beams spurred the search for experimental manifestations of the weak neutral current. Its discovery in 1973 (21, 22) marked an important milestone, as did the observation a decade later (23) of the $W^{\pm}$ (24, 25) and Z^0 (26, 27) bosons. The early years were marked by some inconsistent experimental results and the invention of many alternatives to the $SU(2)_L \otimes U(1)_Y$ theme. How physicists sorted out the correct electroweak theory is a fascinating story but beyond the scope of this review. Instead, I focus on the evidence that now tests and validates the electroweak theory. See Reference 28 for a compact authoritative rendering of the role of precision measurements in establishing the electroweak theory as a law of nature.

SLAC National Accelerator Laboratory (SLAC): located in Menlo Park, California; operated for the U.S. Department of Energy; devoted to particle and accelerator physics, astrophysics, and photon science. See **http://slac.stanford.edu**

European Laboratory for Particle Physics (CERN): straddles the French-Swiss border near Geneva. Its principal research instrument is now the LHC. One of Europe's first common undertakings at its founding in 1954, CERN now includes 20 member states; see **http://cern.ch**

Tevatron: Located at Fermilab, it is the first superconducting synchrotron, operated as a proton-antiproton collider at 2-TeV c.m. energy and luminosity (at the beginning of a run) above $3 \times 10^{32}\ \mathrm{cm}^{-2}\ \mathrm{s}^{-1}$. See **http://www-bd.fnal.gov**

Deutsches Elektronen-Synchrotron (DESY): located in Hamburg, Germany; member of the Helmholtz Association. Develops accelerators and detectors for particle physics and photon science; see **http://zms.desy.de/index_eng.html**

3.1. Tree Level

Following the discovery of neutral-current interactions, the new phenomenon was taken up in a number of νN and νe scattering experiments. Despite their statistical limitations, the neutrino-electron scattering experiments helped guide the convergence to the $SU(2)_L \otimes U(1)_Y$ Standard Model. Under modest universality assumptions, the νe cross-section measurements, combined with measurements of the forward-backward asymmetry that arises from γZ interference in the reaction $e^+e^- \to \mu^+\mu^-$, uniquely selected the $SU(2)_L \otimes U(1)_Y$ chiral couplings of Z to charged leptons (29). Only a short time later, it was reasonable to proclaim that the chiral couplings to all the known quarks and leptons had been uniquely determined in agreement with the $SU(2)_L \otimes U(1)_Y$ theory (30).

Along the way, delicate observations of parity-violating phenomena in atomic physics began to add complementary information. Studies of polarized electron-deuteron scattering (31) confirmed that the neutral-current interactions are parity violating, which also supported the Standard Model.

This impressive progress, punctuated by the discoveries of W and Z, preceded the incisive experiments at the SLAC and CERN Z factories. Measurements of the Z line shape and a determination of the "invisible" width of the Z confirmed the hypothesis that three generations of light neutrinos are present in neutral-current interactions. Not only is the current inference from the invisible width of 2.985 ± 0.009 active light neutrino species (32) consistent with the three observed neutrino species, it leaves little room for decays of Z into exotic weakly interacting particles.

The conclusion that only three active light species exist does not rule out a fourth generation of quarks and leptons, provided that the neutral leptons are heavy enough that their contributions to the invisible width would be negligible—if not zero! A fourth generation is constrained, but not excluded, by what we know of charged-current and neutral-current interactions (33).

Many extensions to the electroweak theory predict the existence of one or more electrically neutral color-singlet Z' gauge bosons (34). The most telling direct searches have been carried out at the Tevatron as searches for direct-channel ($q\bar{q} \to Z'$) resonances in reactions such as $\bar{p}p \to \ell^+\ell^- +$ anything. Translating experimental sensitivity into limits on the mass of a new neutral gauge boson is complicated by the fact that Z' couplings to fermions are model dependent and in some cases even generation dependent. For a representative collection of examples, the Tevatron searches imply that $M_{Z'} \gtrsim 789$ GeV at 95% CL. For a heavy clone of the Standard-Model Z (its only virtue as an example is that it is easy to state), the 95% CL bound is $M_{Z'_{\mathrm{SM}}} > 1030$ GeV (35). Other searches look for evidence of $Z' \to W^+W^-$. Global fits to electroweak parameters and neutral-current studies away from the Z pole are sensitive to a Z'.

The H1 (36) and ZEUS (37–40) experiments at DESY's $e^{\pm}p$ collider HERA compared the momentum-transfer dependence of neutral-current ($e^{\pm}p \to e^{\pm} +$ anything) and charged-current [$e^{\pm}p \to (\bar{\nu}_e, \nu_e) +$ anything] at c.m. energies of 820 and 920 GeV. A recent summary compiled by H1 and ZEUS is given in **Figure 1**. At low values of Q^2, the neutral-current cross section exceeds the charged-current cross section by more than two orders of magnitude

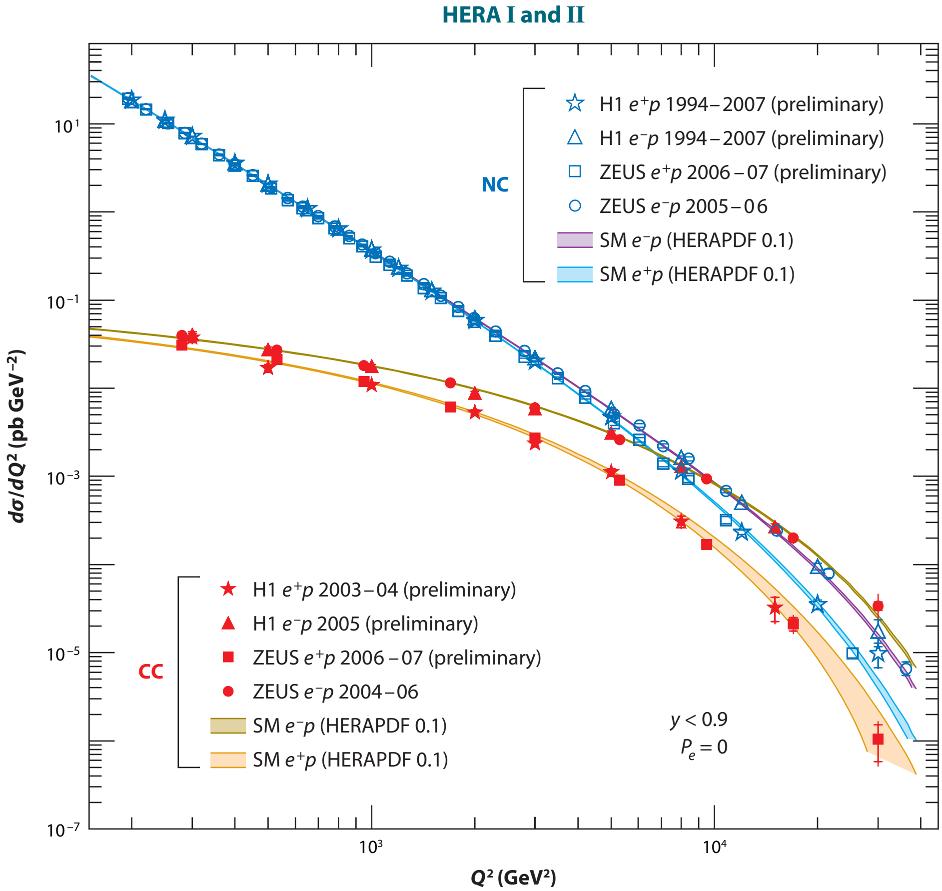


Figure 1

The Q^2 dependence of the neutral-current (NC) and charged-current (CC) cross sections measured by the H1 (36) and ZEUS (37–40) experiments at the HERA $e^{\pm}p$ collider. The curves represent the Standard-Model expectations derived from the HERA parton distribution functions.

because the electromagnetic interaction is much stronger than the weak interaction at long wavelengths. For $Q^2 \gtrsim (M_W^2, M_Z^2)$, the cross sections roughly track each other. This behavior supports the notion that the intrinsic strengths of the weak and electromagnetic interactions are comparable.

The absence of right-handed charged-current interactions is one of the foundational observations on which the $SU(2)_L \otimes U(1)_Y$ theory is built, and it is also a question that has lingered for more than 50 years. Is there a fundamental left-right asymmetry in the laws of nature, or did spontaneous symmetry breaking at some high scale give a large mass to a right-handed gauge boson, creating a low-energy preference for left-handed currents? The second possibility is the vision of left-right symmetric models based on $SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ gauge symmetry (41–43). Searches for right-handed interactions, or for additional $W'^{\pm}$ gauge bosons, are important probes of the electroweak theory (44).

The direct searches at the Tevatron for $W' \to e\nu$, assuming Standard-Model couplings, give a lower bound $M_{W'} > 1000$ GeV at 95% CL (45). A fit to low-energy data bounds the mass of a

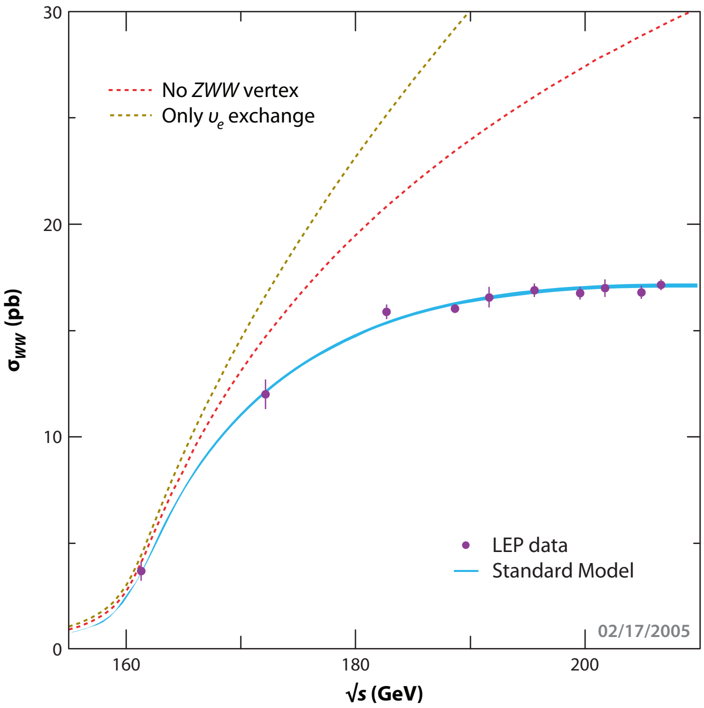


Figure 2

Cross section for the reaction $e^+e^- \to W^+W^-$ measured by the four LEP experiments, together with the full electroweak theory simulation and cross sections that would result from ν exchange alone and from $(\nu + \gamma)$ exchange. If the Z-exchange contribution is omitted (*middle line*) or if both the γ- and Z-exchange contributions are omitted (*upper line*), the calculated cross section grows unacceptably with energy (see **http://lepewwg.web.cern.ch**).

right-handed W_R as $M_{W_R} > 715$ GeV at 90% CL, assuming that its gauge coupling is the same as the $SU(2)_L$ coupling, $g_R = g_L$ (46). Sensitive tests of the Standard Model are ongoing in μ decay (47) and in β decay (48).

LEP: a double ring of synchrotrons providing electron-positron collisions in CERN's 27-km tunnel. Originally constructed as a Z factory, LEP eventually reached a c.m. energy of 209 GeV. See **http://www.cern.ch/Public/en/Research/LEP-en.html**

A noteworthy achievement of the LEP experiments is the validation of the $SU(2)_L \otimes U(1)_Y$ symmetry for the interactions of gauge bosons with fermions and of gauge bosons with gauge bosons in $e^+e^- \to W^+W^-$. This reaction is described by three Feynman diagrams that correspond to s-channel photon and Z^0 exchange and to t-channel neutrino exchange. For the production of longitudinally polarized W bosons, each diagram leads to a $J = 1$ partial-wave amplitude that grows as the square of the c.m. energy, but the gauge symmetry enforces a pattern of cooperation. The contributions of the direct-channel γ- and Z^0-exchange diagrams cancel the leading divergence in the $J = 1$ partial-wave amplitude of the neutrino-exchange diagram. The interplay is shown in **Figure 2**. The measurements compiled by the LEP Electroweak Working Group (see **http://lepewwg.web.cern.ch**) agree well with the benign high-energy behavior predicted by the full electroweak theory.

Tevatron measurements do not directly determine the W^+W^- invariant mass because of the missing energy carried by neutrinos, but they reach beyond the highest energy studied at LEP. The latest contributions, from the D0 (49) and CDF (50) Collaborations, are in agreement with Standard-Model expectations (51, 52), and tighten the bounds on anomalous couplings.

3.2. Flavor-Changing Neutral Currents

Strangeness-changing neutral currents were the object of experimental searches even before the electroweak theory was conceived. It was recognized early on that FCNC effects cannot be isolated in nonleptonic decays. As an example, the transition $s \to d(u\bar{u})$ would be entangled with the charged-current transition $s \to u(d\bar{u})$. Accordingly, decays of hadrons into pairs of leptons have been the favored hunting ground for evidence of FCNC.

The branching fraction $\mathcal{B}(K_L^0 \to \mu^+\mu^-) = (6.84 \pm 0.11) \times 10^{-9}$ (45) closely matches the standard expectation for decay through the (real and virtual) $\gamma\gamma$ intermediate state. The absence of strangeness-changing neutral-current interactions motivated Glashow, Iliopoulos & Maiani (GIM) (53) to advocate adding the charm quark c to the then-familiar u, d, and s, so that quark doublets

$$\begin{pmatrix} u \\ d\cos\theta_C + s\sin\theta_C \end{pmatrix}_L \begin{pmatrix} c \\ s\cos\theta_C - d\sin\theta_C \end{pmatrix}_L , \qquad 20.$$

where θ_C is the Cabibbo angle, would mirror the then-known lepton doublets,

$$\begin{pmatrix} \nu_e \\ e \end{pmatrix}_L \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}_L . \qquad 21.$$

The three-family generalization of the GIM mechanism banishes FCNC at lowest order and greatly suppresses them at loop level (54). Verifying the absence of FCNC therefore tests the structure—and the completeness—of the electroweak theory. The most sensitive experimental search was carried out in the $K^+ \to \pi^+\nu\bar{\nu}$ channel. Brookhaven Experiment 949 has observed three candidates, leading to a branching fraction $\mathcal{B}(K^+ \to \pi^+\nu\bar{\nu}) = 1.73^{+1.15}_{-1.05} \times 10^{-10}$ (55). This rate is consistent, within uncertainties, with the Standard-Model expectation, $\mathcal{B}(K^+ \to \pi^+\nu\bar{\nu}) = (0.85 \pm 0.07) \times 10^{-10}$ (56).

The limits on FCNC involving heavier flavors are less stringent, but they nevertheless raise the question: If new physics is to reveal itself on the 1-TeV scale, why have we seen no sign of FCNC?

Within the Standard Model, the rate anticipated for the decay $D^0 \to \mu^+\mu^-$ is very small: $\mathcal{B}(D^0 \to \mu^+\mu^-) \gtrsim 4 \times 10^{-13}$ (57). The CDF Collaboration bounds $\mathcal{B}(D^0 \to \mu^+\mu^-) < 5.3 \times 10^{-7}$ at 95% CL (58). For a general review of charmed meson decays, see Reference 59. The observation of $D^0 - \bar{D}^0$ mixing (60, 61) has intensified interest in the search for new physics in charmed meson decays. Theoretical expectations are cataloged in References 57 and 62–66.

An informative introduction to FCNC phenomena in B meson decays is given in the *BaBar Physics Book* (67). The current experimental limit on leptonic B_s decays, $\mathcal{B}(B_s \to \mu^+\mu^-) < 5.3 \times 10^{-8}$ at 95% CL (68), approaches Standard-Model sensitivity, $\mathcal{B}(B_s \to \mu^+\mu^-) = (3.6 \pm 0.3) \times 10^{-9}$ (69). The corresponding limit for B^0 is $\mathcal{B}(B_d \to \mu^+\mu^-) < 1.8 \times 10^{-8}$ (68), to be compared with the Standard-Model expectation, $\mathcal{B}(B_d \to \mu^+\mu^-) = (1.1 \pm 0.1) \times 10^{-10}$.

The world sample of top decays remains modest, and consequently the study of rare top decays is less advanced than for K, D, and B mesons. From a search for single top production, the CDF Collaboration reports $\mathcal{B}(t \to ug) < 3.9 \times 10^{-4}$ and $\mathcal{B}(t \to cg) < 5.7 \times 10^{-3}$ at 95% CL (70), improving earlier limits from LEP. The latter result is to be compared with the Standard-Model expectation, $\mathcal{B}(t \to cg) \approx 10^{-10}$ (71). A study of top pair production yields $\mathcal{B}(t \to Zc) < 37\%$ at 95% CL (72).

In charm and top decays, plenty of room remains to search for physics beyond the Standard Model, as experiments approach Standard-Model sensitivity. However, the absence of FCNC at tree level is firmly established. What we already know about (the suppression of) FCNC

D0: one of two experiments investigating proton-antiproton collisions at the Tevatron; named for its position on the Tevatron ring. See **http://www-d0.fnal.gov**

Collider Detector at Fermilab (CDF): one of two experiments investigating proton-antiproton collisions at the Tevatron; see **http://www-cdf.fnal.gov**

GIM mechanism: observation by Glashow, Iliopoulos & Maiani (53) that FCNC interactions vanish at tree level and are strongly inhibited at higher orders, provided that quarks (and leptons) occur in $SU(2)_L$ doublets; argued for the necessity of the charm quark phenomena both challenges and provides opportunities to uncover many varieties of physics beyond the Standard Model, including dynamical electroweak symmetry breaking (73–75) and supersymmetry without auxiliary conditions (76). The existing constraints have stimulated conjectures about so-called minimal flavor violation (77) and approximate generational symmetries (78).

The search for FCNC effects in heavy quark decays is an example of how high-sensitivity studies at low energies can complement direct discovery physics at the LHC. Experimental searches for lepton-flavor violation offer another window on new physics in the neutral-current sector (79, 80).

3.3. Tests of the CKM Paradigm

A generation ago, the Cabibbo hypothesis (11) brought clarity to a wealth of information about semileptonic decays of mesons and hyperons. Transcribed to modern language, the charged-current interactions among the light quarks are specified by

$$\mathcal{L}_{\mathrm{CC}}^{(q)} = -\frac{g}{\sqrt{2}}\bar{u}_{\mathrm{L}}\gamma^{\mu}d_{\theta\mathrm{L}}W_{\mu}^{+} + \mathrm{h.c.}, \tag{22}$$

where g is the $SU(2)_{\mathrm{L}}$ gauge coupling and

$$d_{\theta} = d\cos\theta_{\mathrm{C}} + s\sin\theta_{\mathrm{C}}. \tag{23}$$

Equation 22 matches the charged-current interaction among leptons,

$$\mathcal{L}_{\mathrm{CC}}^{(\ell)} = -\frac{g}{\sqrt{2}}\bar{e}_{\mathrm{L}}\gamma^{\mu}\nu_{\mathrm{L}}W_{\mu}^{-} + \mathrm{h.c.}, \tag{24}$$

and so expresses the universality of the charged-current weak interactions. Tests of the Cabibbo universality hypothesis relating the strengths of $u \leftrightarrow d$, $u \leftrightarrow s$, and $\nu \leftrightarrow e$ transitions are reviewed in Reference 81.

In a prescient paper, following the GIM (53) call for a fourth quark that would be the charged-current partner of the orthogonal combination $s_{\theta} = s\cos\theta_{\mathrm{C}} - d\sin\theta_{\mathrm{C}}$ but before the discovery of charm, Kobayashi & Maskawa (12) generalized Cabibbo's hypothesis to three quark generations in order to accommodate *CP* violation. Quark mixing is expressed by the 3×3 unitary matrix defined in Equation 5, known as the CKM matrix. The authors' key insight is that an $n \times n$ unitary matrix can be parameterized in terms of $n(n-1)/2$ real mixing angles and $(n-1)(n-2)/2$ complex phases, after the freedom to redefine the phases of quark fields has been taken into account. The phase angle present in the 3×3 case could, they suggested, account for *CP* violation. This simple conjecture has far-reaching implications (82, 83). We now know of three generations of leptons (Equation 1) and quarks (Equation 3)—a good beginning.

A simple test for the completeness of the CKM picture is to determine whether the magnitudes $|V_{ij}|$ are consistent with the hypothesis that the matrix is unitary. Researchers have devoted particular attention to the first row of the CKM matrix, looking for deviations from the unitarity requirement

$$\mathcal{S}_u \equiv |V_{ud}|^2 + |V_{us}|^2 + |V_{ub}|^2 = 1, \tag{25}$$

which would signal new physics. (Because $|V_{ub}|^2 \ll 1$, this is essentially a test of the Cabibbo picture.) For several years, the sum $\mathcal{S}_u$ lay a couple of standard deviations below unity. Recent kaon decay studies have raised the value of $|V_{us}|$, so that $\mathcal{S}_u = 0.9999 \pm 0.0010$ (84). Ongoing studies of neutron decays should resolve a persistent lifetime puzzle (85) and may lead to an improved determination of $|V_{ud}|$.

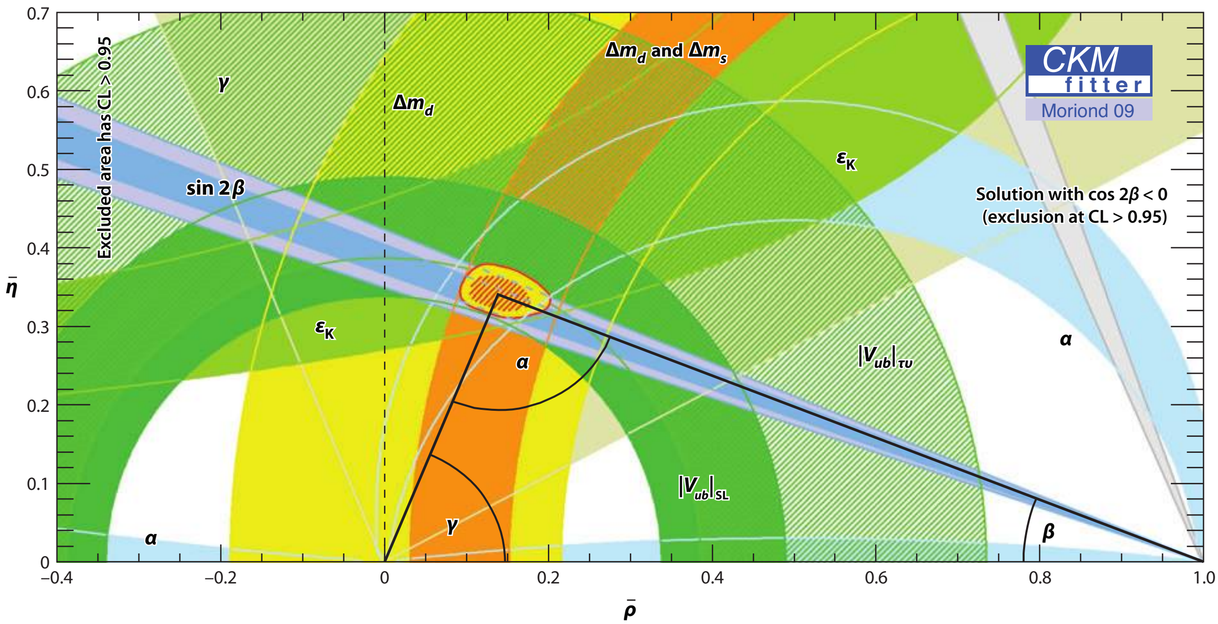


Figure 3

Constraints in the ($\bar{\rho}$, $\bar{\eta}$) plane as of March, 2009. The orange hashed region shows the global combination at 68% CL (90).

Immense experimental effort has produced a rich library of information about decays (both common and rare), neutral-particle mixings, and *CP* violation (in *K* and *B* decays) (86–88). One application of that body of knowledge has been to probe in depth the unitarity of the CKM matrix $\mathsf{V}\mathsf{V}^\dagger = \mathsf{I}$, where I is the 3×3 identity, through examination of $\sum_i V_{ij}V^*_{ik} = \delta_{jk}$ and $\sum_i V_{ij}V^*_{kj} = \delta_{ik}$. The six vanishing conditions may be represented as triangles in the complex plane, each with an area proportional to $\mathrm{Im}[V_{ij}V_{k\ell}V^*_{i\ell}V^*_{kj}]$, a parameterization-independent measure of *CP* violation (89). Comprehensive analyses have been carried out over a number of years by the CKMfitter (90) and UT*fit* (91) Collaborations.

The most commonly displayed unitarity triangle, shown in **Figure 3**, is constructed from the constraint

$$V_{ud}V^*_{ub} + V_{cd}V^*_{cb} + V_{td}V^*_{tb} = 0. \qquad 26.$$

It is conventional to normalize the triangle, dividing the complex vector for each leg by the well-determined $V_{cd}V^*_{cb}$. The vertices of the triangle are then (0, 0), (1, 0), and ($\bar{\rho}$, $\bar{\eta}$). Among the tests available in this formalism are whether the triangle closes and whether different data sets yield a common vertex, ($\bar{\rho}$, $\bar{\eta}$). The plot in **Figure 3**, which is representative of recent work, shows consistency among many experimental constraints. That the imaginary coordinate $\bar{\eta}$ differs from zero shows that the Kobayashi–Maskawa mechanism is at work. A crucial prediction, that *CP* violation in *K* physics is small because of flavor suppression but *CP* violation should be appreciable in *B* physics, is fulfilled. More detailed analysis shows that the Kobayashi–Maskawa mechanism is the dominant source of *CP* violation in meson decays. As discussed in Section 3.2, new physics contributions are extremely small in $s \leftrightarrow d$, $b \leftrightarrow d$, $s \leftrightarrow b$, and $c \leftrightarrow u$ transitions. For summaries of tests of the CKM paradigm in flavor physics, see Reference 92 for an experimental perspective and a look ahead and Reference 69 for a theoretical perspective.

A global fit (93), within the framework of the three-generation Standard Model, yields the following magnitudes $|V_{ij}|$ for the CKM matrix elements:

$$\begin{pmatrix} 0.97419 \pm 0.00022 & 0.2257 \pm 0.0010 & 0.00359 \pm 0.00016 \\ 0.2256 \pm 0.0010 & 0.97334 \pm 0.00023 & 0.0415^{+0.0010}_{-0.0011} \\ 0.00874^{+0.00026}_{-0.00037} & 0.0407 \pm 0.0010 & 0.999133^{+0.000044}_{-0.000043} \end{pmatrix}. \qquad 27.$$

The consistency of the CKM picture does not yet exclude a fourth generation of quarks. Direct constraints on $|V_{tb}|$ are consistent with a value near unity, but they are not yet terribly restrictive. Global fits to the precision electroweak data allow mixing between the third and fourth families at the level seen between the first and second families (94).

Finally, the robustness of the CKM unitarity triangle does not mean that there is no new physics to be found. The unitarity triangle analysis is mainly sensitive to processes that change flavor by two units. Even in the well-studied rare K and B decays (processes that change flavor by one unit), many examples of new physics that could have passed the unitarity-triangle screen—supersymmetry, little Higgs models with T-parity, and warped extra dimensions—could give large departures (95). New sources of CP violation and FCNC occur in models that do not enforce minimal flavor violation. As described in Section 3.2, there is ample space between current bounds and Standard-Model expectations in many rare decays. Because the unitarity triangle is described well by the Standard Model, it will be useful to examine CP violation in $b \to s$ transitions and rare decays, where Standard-Model contributions are small. One specific scenario, involving extra $U(1)'$ interactions, is presented in Reference 96, and a claimed sign of new physics in $b \to s$ transitions is given in Reference 97.

The ability of the electroweak theory incorporating CKM mixing to account for—and predict—a vast number of observables in flavor physics is highly impressive. We must remember, however, that experiments have validated a framework, not an explanation. Just as the Standard Model makes no predictions for quark and lepton masses, it has nothing to say about the mixing angles and the Kobayashi–Maskawa phase. These can arise in the electroweak theory, but we do not know how. If quark and lepton masses and mixings are indeed generated by the Higgs mechanism, then—in the words of Veltman (98)—the Higgs boson must know something we do not know.

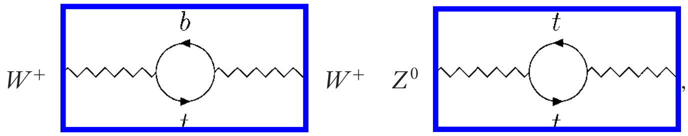
3.4. Loop Level

We have just recalled some of the ways in which experiment has tested the consequences of the spontaneously broken $SU(2)_L \otimes U(1)_Y$ gauge theory of the electroweak interactions and in which it has probed with increasing acuity the inferences from earlier experiments on which the electroweak theory was founded. The major predictions for electroweak phenomenology have been confirmed—among them, the existence of neutral-current interactions, the existence and mass scale of the $W^{\pm}$ and Z^0, and the need for the charm quark. The idealizations that shaped the structure of the theory—including the absence of right-handed charged currents and the absence of FCNC interactions—have proved to be exceptionally robust. Only the idealization that the neutrinos are massless has required revision, and that is for many purposes an inessential change.

The electroweak theory is a quantum field theory. Once the elementary interactions have been set by hypothesis or by experimental determinations, we have the opportunity to compute quantum corrections to observables and subject the theory to precise experimental tests. An accessible introduction to the basic techniques can be found in Reference 99. The program is straightforward in principle but very demanding in practice. The mounting precision of experiments has inspired waves of detailed theoretical calculations that are heroic in proportion (100).

If all the parameters of a theory are known (and if the theory is presumed complete), then a measured observable may be compared with the calculated value to test the theory. The electroweak theory has been a work in progress over the period in which precise measurements became available because several key parameters have been unknown.

Before the top quark was discovered in 1995, quantum corrections to electroweak observables gave indications that the weak-isospin partner of b would be much more massive than the other quarks. For example, the quantum corrections to the Standard-Model predictions for M_W (Equation 18) and for M_Z (Equation 19) arise from different quark loops:



$t\bar{b}$ for M_W and $t\bar{t}$ (or $b\bar{b}$) for M_Z. These quantum corrections alter the link between the W and Z boson masses, so that

$$M_W^2 = M_Z^2\left(1 - \sin^2\theta_{\mathrm{W}}\right)(1 + \Delta\rho), \qquad 28.$$

where

$$\Delta\rho \approx \Delta\rho^{(\mathrm{quarks})} = \frac{3G_{\mathrm{F}}m_t^2}{8\pi^2\sqrt{2}}. \qquad 29.$$

The strong dependency on m_t^2 is characteristic and accounts for the sensitivity of electroweak observables to the top quark mass.

If all other parameters were known, to estimate m_t one could choose for any measurement the value of m_t that gave the closest agreement between calculation and experiment, test for consistency among various measurements, and average over different observables. In practice, the global fits allow for variations in a number of parameters. The top mass favored by simultaneous fits to many electroweak observables is shown as a function of time in **Figure 4**. By the end of 1994, the indirect determinations favored $m_t \approx (175 \pm 25)$ GeV, successfully anticipating the masses reported in the discovery papers: $176 \pm 8 \pm 10$ GeV for CDF and $199^{+19}_{-21} \pm 22$ GeV for D0. Today, direct measurements at the Tevatron determine the top quark mass to a precision of 0.75%, $m_t = (173.1 \pm 1.3)$ GeV (101), far more precise than the indirect determinations.

Measurements on and near the Z^0 pole by the LEP experiments ALEPH, DELPHI, L3, and OPAL (104) and by the SLD experiment at the Stanford Linear Collider (105) were decisive in testing and refining the electroweak theory (106). Global analysis projects that have been distinguished for their thoroughness and continuity include the LEP Electroweak Working Group (102; also see **http://lepewwg.web.cern.ch**), incorporating the ZFITTER (107, 108) and TOPAZ0 (109, 110) codes, and the Particle Data Group (32). These projects were recently joined by the Tevatron Electroweak Working Group (see **http://tevewwg.fnal.gov**) and the Gfitter initiative (103).

What has been achieved overall is a comprehensive test of the electroweak theory, as a quantum field theory, at a precision of one part in a thousand for several observables. A representative comparison of best-fit calculations with observations is shown in **Figure 5** (103), which displays for each observable the difference between fitted and measured values, weighted by the inverse of the experimental standard deviation. (See **http://lepewwg.web.cern.ch** and Reference 102 for the corresponding information from the LEP Electroweak Working Group and Reference 32 for the Particle Data Group's version.) For only 1 observable out of 20—the forward-backward asymmetry in the reaction $e^+e^- \rightarrow b\bar{b}$ on the Z resonance—does the difference exceed two standard deviations. The global fits yield excellent determinations of Standard-Model parameters, including the weak mixing parameter.

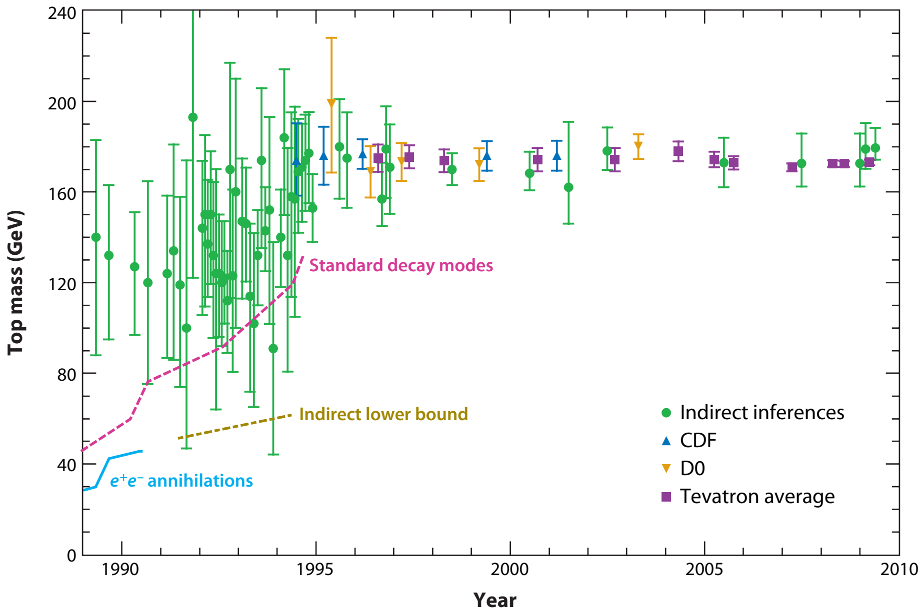


Figure 4

Indirect determinations of the top quark mass from fits to electroweak observables (*green circles*) and 95%-CL lower bounds on the top quark mass inferred from direct searches in e^+e^- annihilations (*solid light blue line*) and in $\bar{p}p$ collisions, assuming that standard decay modes dominate (*broken magenta line*). An indirect lower bound, derived from the W boson width inferred from $\bar{p}p \to (W \text{ or } Z) + \text{anything}$, is shown as the dotted-dashed brown line. Selected direct measurements of m_t by the CDF (*dark blue triangles*) and D0 (*inverted orange triangles*) Collaborations are plotted. The Tevatron average from direct observations is shown as magenta squares. The most recent indirect determinations are from References 102 and 103 and from **http://lepewwg.web.cern.ch**. The evolution of knowledge of m_t may be traced through the current *Review of Particle Physics* (45) and previous editions.

3.5. Evidence for Higgs Boson Interactions

An important asset of global fits to many observables is their sensitivity to virtual effects and, thus, to parameters that have not been measured directly. The successful inference of the range of top quark masses is a prime example. Now that m_t is measured at high precision, it becomes a fixed parameter in the global fits, which may probe for the next unknown quantity.

Figure 6 shows how the goodness of the LEP Electroweak Working Group's Winter 2009 global fit depends upon M_H. The fit is evidently improved by the inclusion of quantum corrections involving a Higgs boson that has Standard-Model interactions with the electroweak gauge bosons $W^{\pm}$ and Z. A satisfactory fit does not prove that the Standard-Model Higgs boson exists, but it does offer guidance for the search and sets up a consistency check when a putative Higgs boson is observed. The inferred range is consistent with the conditional upper bound, $M_H \lesssim 1$ TeV, derived in Section 4.1. It is important to note that, whereas the global fits give evidence for the effect of the Higgs boson in the vacuum, they do not have any sensitivity to couplings of the Higgs boson to fermions free of the assumption that Higgs–Yukawa couplings set the fermion masses.

The precision electroweak measurements on their own argue for $M_H \lesssim 163$ GeV, a one-sided 95%-CL limit derived from $\Delta\chi^2 = 2.7$ for the blue band in **Figure 6**. Imposing the exclusion

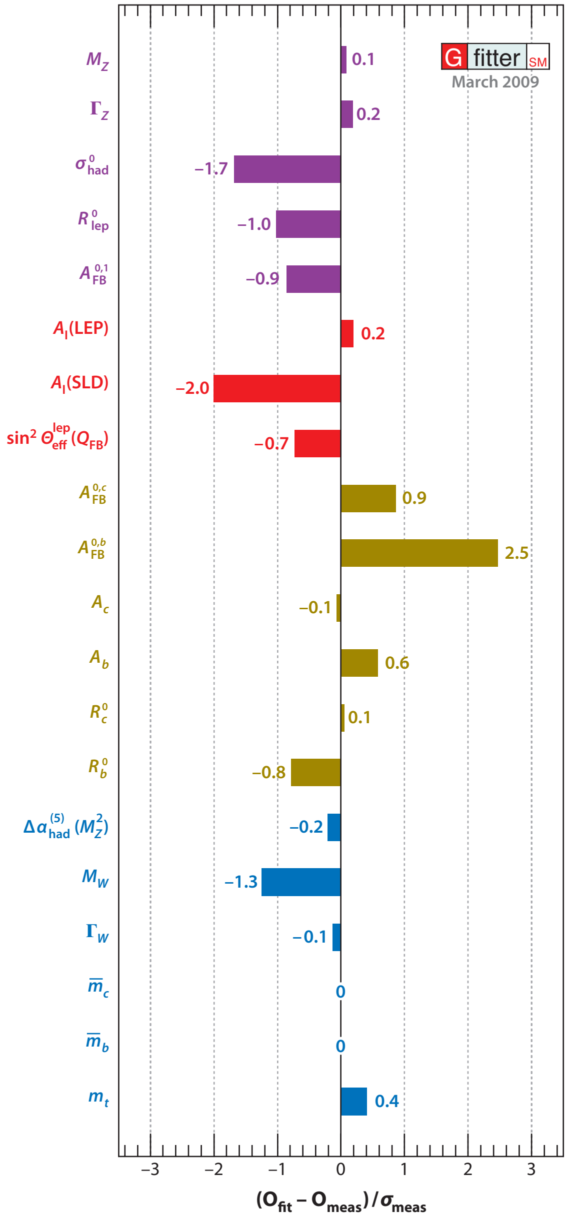


Figure 5

Pull values comparing Gfitter complete fit results with experimental determinations (103).

$M_H > 114.4$ GeV from the LEP searches leads to an upper bound of $M_H \lesssim 191$ GeV. The Particle Data Group (32) and Gfitter (103) analyses lead to similar conclusions.

The Higgs boson masses, favored by the global fits of the LEP Electroweak Working Group, $M_H = 90^{+36}_{-27}$ GeV (see **http://lepewwg.web.cern.ch**), Gfitter, 83^{+30}_{-23} GeV (103), or the Particle Data Group, 70^{+28}_{-22} GeV (32), lie in the region excluded by direct searches at LEP. Chanowitz (113, 114) has cautioned that the values of M_H preferred by fits to different observables are not entirely

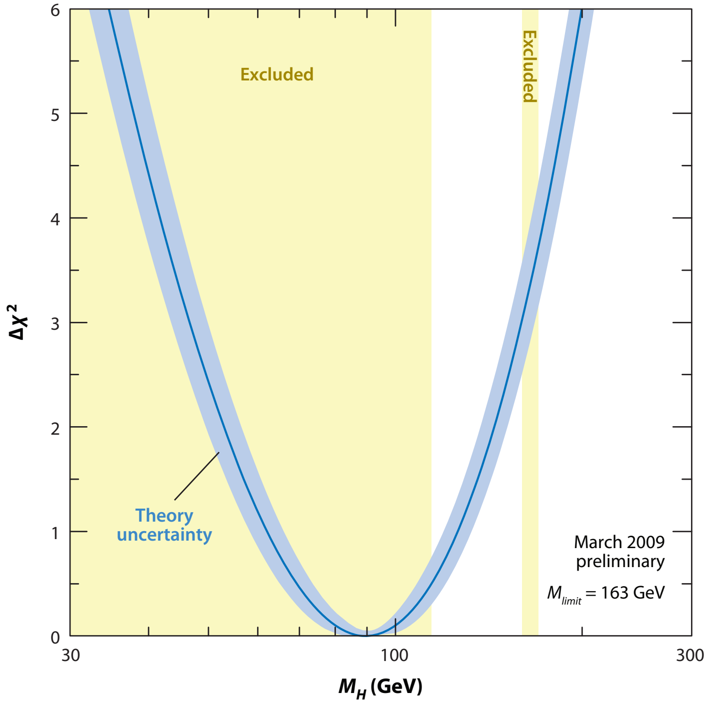


Figure 6

$\Delta\chi^2 = \chi^2 - \chi^2_{\mathrm{min}}$ from a fit to an ensemble of electroweak measurements as a function of the Standard-Model Higgs boson mass. The solid blue line is the result of the fit. The blue band represents an estimate of the theoretical uncertainty due to missing higher-order corrections. The regions shaded in yellow denote the 95%-CL lower bound on $M_H > 114.4$ GeV from direct searches at LEP (111) and the Tevatron exclusion at 95% CL between 160 and 170 GeV (112). Analysis is by the LEP Electroweak Working Group (see **http://lepewwg.web.cern.ch**).

consistent. The scatter is illustrated in the case of the Gfitter analysis in **Figure 7**. In particular, the forward-backward asymmetry in $e^+e^- \to b\bar{b}$ on the Z resonance ($A^{0,b}_{\mathrm{FB}}$) is best reproduced with $M_H \approx 400$ GeV. This observable is the most discrepant,[1] at $\gtrsim 2.5\sigma$, with the overall fits (cf. **Figure 5**). Omitting it (on the hypothesis that it is particularly sensitive to new physics) would improve the global fits but lead to a small Higgs boson mass that would coexist uncomfortably with the LEP exclusion: The Gfitter best-fit range would move to 61^{+30}_{-26} GeV. Whether the spread of Higgs boson masses preferred by different sensitive observables points to physics beyond the Standard Model or represents insignificant scatter is a tantalizing question.

3.6. The Weak Mixing Parameter at Low Scales

The extraordinary precision of measurements on the Z^0 pole has given them a decisive weight in our assessment of the electroweak theory. They are, however, blind to new physics that does not directly modify the Z^0 properties. A heavy Z' that does not mix appreciably with Z^0 is an important

[1]For the purpose of this discussion, I set aside the anomalous magnetic moment of the muon (115), for which the Standard-Model prediction remains somewhat uncertain. See Reference 116.

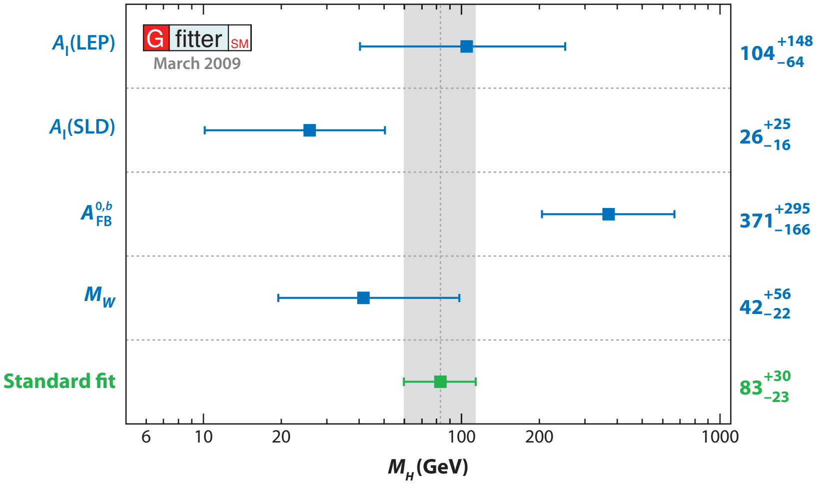


Figure 7

Determination of the Higgs boson mass excluding all the sensitive observables from the Gfitter standard fit, except for the one given (103).

example. For this reason, experiments off the Z^0 pole, even of lower precision, command our attention—particularly in the search for physics beyond the Standard Model.

The weak mixing parameter is defined in terms of (running) couplings,

$$\sin^2\theta_W(Q) = \frac{\alpha(Q)}{\alpha_2(Q)} = \frac{1/\alpha_2(Q)}{1/\alpha_Y(Q) + 1/\alpha_2(Q)}, \qquad 30.$$

so its value depends on the scale at which it is measured. A familiar illustration occurs in unified theories of the strong, weak, and electromagnetic interactions, which predict the value of the weak mixing parameter at low scales. The prototype is the $SU(5)$ unified theory (117): At the unification scale $U \approx 10^{15}$ GeV, the running couplings are simply related:

$$\left.\begin{aligned} 1/\alpha_2 &= 1/\alpha_U \\ 1/\alpha_Y &= \tfrac{5}{3}\cdot 1/\alpha_U \\ 1/\alpha &= \tfrac{8}{3}\cdot 1/\alpha_U \end{aligned}\right\}, \qquad 31.$$

where α_U is the common value of the $SU(3)_c$, $SU(2)_L$, and $U(1)$ couplings. At the $SU(5)$ unification scale, $\sin^2\theta_W(U) = \frac{3}{8}$. How does $\sin^2\theta_W$ evolve? In leading logarithmic approximation and at high scales (118),

$$\sin^2\theta_W(Q) = \frac{3}{8} - \frac{5}{8}(b_1 - b_2)\alpha(Q)\log\left(Q^2/U^2\right), \qquad 32.$$

where the beta functions $4\pi b_1 = -4n_g/3 - n_H/10$ and $4\pi b_2 = (22 - 4n_g)/3 - n_H/6$ determine the evolution of $1/\alpha_1$ and $1/\alpha_2$. Here n_g is the number of fermion generations, and n_H is the number of Higgs doublets. The weak mixing parameter decreases as Q decreases from the unification scale U. At the Z boson mass, $\sin^2\theta_W(M_Z)\big|_{SU(5)} \approx 0.21$, near (but not near enough) to the measured value, $\sin^2\theta_W(M_Z)\big|_{\text{exp}} = 0.23119 \pm 0.00014$ in the $\overline{\text{MS}}$ scheme (32).

In the range of scales directly accessible to experiment, the evolution of the weak mixing parameter is predicted within the electroweak theory itself. The expectations of a higher-order renormalization-group analysis (119) are depicted in **Figure 8**. A detailed comparison with experiment is given in Reference 32. Here are some of the main points. The parity-violating left-right

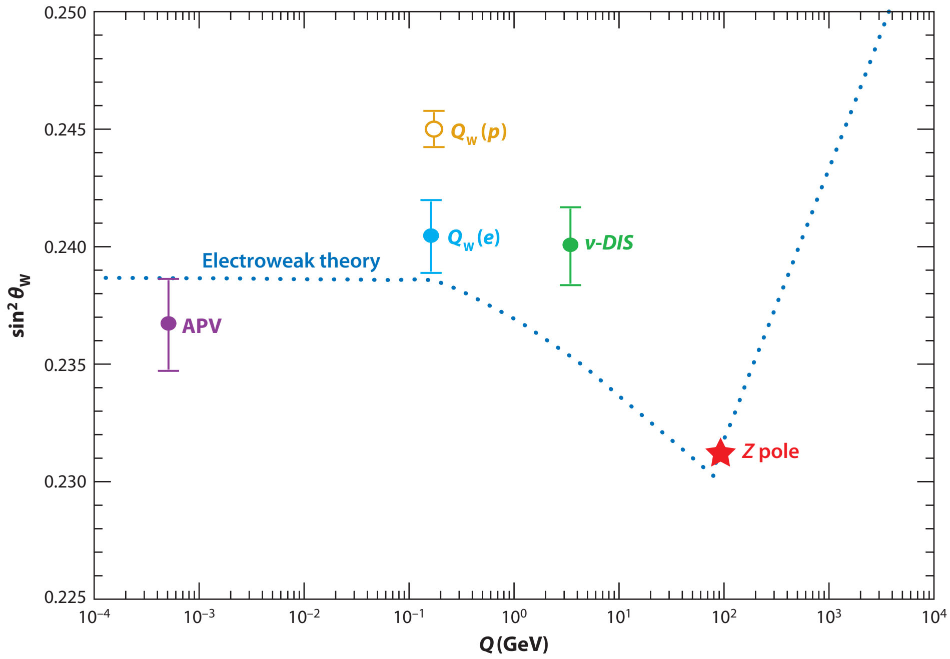


Figure 8

Evolution of the weak mixing parameter $\sin^2 \theta_W$ in the $\overline{\mathrm{MS}}$ scheme (*dark blue dotted curve*) (119). The minimum occurs at $Q = M_W$, where the β function for the weak mixing parameter changes sign as the influence of weak boson loops drops out. The selected data are from atomic parity violation (APV; *purple*) (120), Møller scattering [$Q_W(e)$; *light blue*] (121), and deep-inelastic νN scattering (DIS; *green*) (122, 123). Also indicated (*orange open circle*) is the uncertainty projected for the Q_{weak} experiment (see **http://www.jlab.org/qweak**).

asymmetry observed (121) in polarized Møller scattering, $e^-e^- \to e^-e^-$, at SLAC establishes the low-energy running of $\sin^2\theta_W$ at more than six standard deviations, and it is in reasonable agreement with the prediction at $Q^2 = 0.026\ \text{GeV}^2$. After important improvements in the connection between the measured quantity and $\sin^2\theta_W$, the most telling measurement of atomic parity violation (120) agrees with the electroweak theory within approximately one standard deviation. The Q_{weak} experiment (see **http://www.jlab.org/qweak**), to be mounted at Jefferson Laboratory at the beginning of 2010, aims for a 0.3% determination of $\sin^2 \theta_W$ in parity-violating scattering of polarized electrons on protons at $Q^2 = 0.03\ \text{GeV}^2$.

The NuTeV experiment at Fermilab determined $\sin^2 \theta_W$ by measuring neutral-current and charged-current cross sections for deep-inelastic νN and $\bar{\nu} N$ scattering (122, 123). NuTeV's result, which lies some three standard deviations above the electroweak theory expectation, has been subjected to intense scrutiny. For the moment, enough ambiguity attends the dependency on fine details of parton distribution functions, the influence of nuclear targets, and various isospin-violating effects that the significance of the NuTeV anomaly is subject to debate. A catalog of some new physics interpretations is given in Reference 124. Many of these [e.g., new Z' gauge bosons (125), leptoquarks, etc.] can be tested at the LHC. New low-energy experiments can test the NuTeV measurement and constrain interpretations. The NuSOnG concept put forward for the Tevatron (126) would supplement deep-inelastic νN scattering with high-statistics measurements of νe and $\bar{\nu} e$ elastic scattering to test for new physics (127, 128) in the neutrino sector.

Fermi National Accelerator Laboratory (Fermilab): located in Batavia, Illinois; operated for the U.S. Department of Energy; home to the Tevatron. See **http://www.fnal.gov**

The LEP-2 measurements at energies between the Z pole and the top energy of 209 GeV were broadly in agreement with Standard-Model expectations (129). Measurements by the CDF (130, 131) and D0 (132) experiments of the forward-backward asymmetry in the reaction $\bar{p}p \to (Z, \gamma^*) +$ anything $\to e^+e^- +$ anything agree with leading-order predictions in the Standard Model over the range of invariant masses $50\,\text{GeV} \lesssim \mathcal{M}(e^+e^-) \lesssim$ few hundred gigaelectronvolts. With the 10 fb^{-1} of data expected by the end of Run II, a measurement of the running of $\sin^2\theta_W$ at an interesting level of precision could be achieved before the LHC experiments pronounce on this subject. For a prospectus of low-energy tests of the weak interaction, see Reference 133.

3.7. The Scale of Fermion Mass Generation

It is no exaggeration to say that the origin of the quark and lepton masses is shrouded in mystery. Within the standard electroweak theory, the overall scale of the fermion masses is set by the vacuum expectation value $v/\sqrt{2} \approx 174$ GeV of the Higgs field, but each fermion mass $m_i = \zeta_i v/\sqrt{2}$ involves a distinct Yukawa coupling ζ_i, as we saw in Equation 16. The Yukawa couplings that reproduce the observed quark and charged-lepton masses range over many orders of magnitude, from $\zeta_e \approx 3 \times 10^{-6}$ for the electron to $\zeta_t \approx 1$ for the top quark, as shown in **Figure 9**. Their origin is unknown. In an important sense, therefore, all fermion masses involve physics beyond the Standard Model.

Although the electroweak theory shows how fermion masses may arise, we cannot be sure that finding the Higgs boson, or understanding electroweak symmetry breaking, will clarify the origin

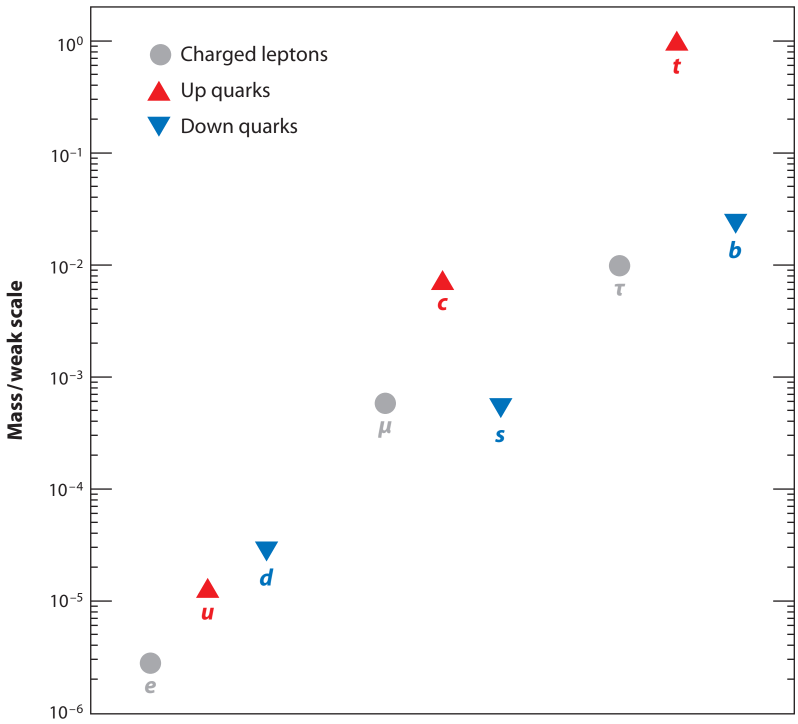


Figure 9

Yukawa couplings $\zeta_i = m_i/(v/\sqrt{2})$ inferred from the masses of the quarks and charged leptons (45).

of fermion masses. This is because we do not know that fermion masses are set on the electroweak scale. This point merits closer examination.

Hard (and soft) fermion mass: A hard mass term is one introduced directly into the Lagrangian of a theory. It persists to arbitrarily high energy scales and, in the case of the electroweak theory, would lead to uncontrolled growth of amplitudes at high energies. A soft mass term that arises through spontaneous breaking of a gauge symmetry does not persist above the scale of symmetry breaking

Effective field theory: a description valid over a particular range of energies or distance scales, based on the degrees of freedom most relevant to the phenomena that occur there. Nonlocal interactions mediated by virtual heavy particles are replaced by local interactions that yield the same low-energy limit. The effective theory can only be a valid description of physics at energies below the masses of the heavy particles and must be superseded by a more complete (but perhaps still effective theory) on that energy scale

The observation of a nonzero fermion mass ($m_i \neq 0$) implies that the electroweak gauge symmetry $SU(2)_L \otimes U(1)_Y$ is broken (cf. Section 2), but electroweak symmetry breaking is only a necessary, not a sufficient, condition for the generation of fermion mass. In the Standard-Model framework, some new physics (at an unknown scale) must give rise to the Yukawa couplings. The logical division of labor between a mechanism for electroweak symmetry breaking and an origin of fermion masses is made explicit in the simple technicolor models (134, 135), discussed in Section 4.4.3 below. In the sparest versions of such models, electroweak symmetry breaking is driven by a gauge interaction that becomes strongly coupled on the electroweak scale. The gauge bosons acquire masses, but the fermions remain massless. So-called extended technicolor models (136–138) invoke additional interactions at a much higher scale (of order 100 TeV) to explain the light quark masses.

Within the framework of the $SU(2)_L \otimes U(1)_Y$ gauge theory, partial-wave unitarity sets a model-independent upper bound on the energy scale of fermion mass generation (139). The strategy is simply to add explicit fermion mass terms to the electroweak Lagrangian, rather than the Yukawa terms of Equation 16. Explicit Dirac mass terms link the left-handed and right-handed fermions, thus violating the $SU(2)_L \otimes U(1)_Y$ gauge symmetry of the electroweak theory. If they persist to arbitrarily high energies, such hard masses destroy the renormalizability of the theory. However, it may be overly ambitious to demand that a theory make sense at all energies. Accordingly, we consider the explicit fermion masses in the framework of an effective field theory valid over a finite range of energies, to be supplanted at higher energies by a theory that entails a different set of degrees of freedom (140).

Because the gauge symmetry is broken in a theory with explicit fermion masses m_i, at lowest order in perturbation theory, scattering amplitudes for the production of pairs of longitudinally polarized gauge bosons in fermion-antifermion annihilations grow with c.m. energy roughly as $G_F m_i E_{cm}$. (In the standard electroweak theory, this behavior is canceled by the contribution of direct-channel Higgs-boson exchange.) The resulting partial-wave amplitudes saturate partial-wave unitarity for the Standard Model with a Higgs mechanism at a critical c.m. energy (139, 141, 142),

$$\sqrt{s_i} \simeq \frac{4\pi\sqrt{2}}{\sqrt{3\eta_i} G_F m_i} = \frac{8\pi v^2}{\sqrt{3\eta_i} m_i}, \qquad 33.$$

where $\eta_i = 1(3)$ for leptons (quarks). As usual, the parameter v sets the scale of electroweak symmetry breaking. If the electron mass were hard, the critical energy would be $\sqrt{s_e} \approx 1.7 \times 10^9$ GeV; the corresponding energy for the top quark would be $\sqrt{s_t} \approx 3$ TeV. The fact that a hard electron mass would only imply a saturation of partial-wave unitarity at a prodigiously high energy means that, although the behavior of $\sigma(e^+e^- \to W^+W^-)$ shown in **Figure 2** validates the gauge symmetry of the electroweak theory, it does not establish that the theory is renormalizable (139, 143).

4. THE AGENT OF ELECTROWEAK SYMMETRY BREAKING

4.1. The Significance of the 1-TeV Scale

The electroweak theory does not give a precise prediction for the mass of the Higgs boson, but a thought experiment leads through a unitarity argument (144) to a conditional upper bound on the Higgs boson mass that sets a key target for experiment.

Consider two-body collisions among $W^\pm$, Z^0, and H. It is straightforward to compute the scattering amplitudes $\mathcal{M}$ at high energies and to make a partial-wave decomposition, according to $\mathcal{M}(s,t) = 16\pi \sum_J (2J+1) a_J(s) P_J(\cos\theta)$. Most channels "decouple," in the sense that partial-wave amplitudes are small at all energies (except very near particle poles or at exponentially large energies) for any value of the Higgs boson mass M_H. Of interest are four neutral channels,

$$W_0^+ W_0^-, \quad \frac{Z_0 Z_0}{\sqrt{2}}, \quad \frac{HH}{\sqrt{2}}, \quad \text{and} \quad HZ_0, \tag{34.}$$

where the subscript 0 denotes the longitudinal polarization states and where the factors of $\sqrt{2}$ account for identical particle statistics. For these, the s-wave amplitudes are all asymptotically constant (i.e., well behaved) and proportional to $G_F M_H^2$. In the high-energy limit $s \gg M_H^2, M_W^2, M_Z^2$,

$$(a_0) \to \frac{-G_F M_H^2}{4\pi\sqrt{2}} \cdot \begin{bmatrix} 1 & 1/\sqrt{8} & 1/\sqrt{8} & 0 \\ 1/\sqrt{8} & 3/4 & 1/4 & 0 \\ 1/\sqrt{8} & 1/4 & 3/4 & 0 \\ 0 & 0 & 0 & 1/2 \end{bmatrix}. \tag{35.}$$

Requiring that the largest eigenvalue respect the partial-wave unitarity condition $|a_0| \leq 1$ yields

$$M_H \leq \left(\frac{8\pi\sqrt{2}}{3G_F}\right)^{1/2} \approx 1\,\text{TeV} \tag{36.}$$

as a condition for perturbative unitarity.

If the Higgs boson mass respects the bound shown in Equation 36, weak interactions remain weak at all energies, and perturbation theory is everywhere reliable. If the Higgs boson mass exceeds 1 TeV, perturbation theory breaks down, as weak interactions among $W^\pm$, Z, and H become strong on the 1-TeV scale. This means that (within the Standard Model) the features familiar in strong-interaction physics at energies near 1 GeV would characterize electroweak boson interactions at energies near 1 TeV. More generally, the implication is that something new—a Higgs boson, strong scattering, or other new physics—is to be found in electroweak interactions at energies not much larger than 1 TeV.

Tighter constraints—in the form of upper and lower bounds on the mass of the Higgs boson—follow from the demand that the electroweak theory be a consistent (and complete) quantum field theory up to a specified energy scale Λ.[2] For a light Higgs boson, the $t\bar{t}H$ Yukawa coupling introduces quantum corrections that may destabilize the Higgs potential (Equation 13) so that the electroweak vacuum state characterized by Equation 17 is no longer the state of minimum energy. The perturbative analysis is explained carefully in Reference 145. For a specified value of the top quark mass, the requirement that the broken-symmetry vacuum of the electroweak theory be the absolute minimum of the (radiatively corrected) Higgs potential gives a lower bound on the Higgs boson mass. For a cutoff $\Lambda = 1$ TeV, the lower bound is (146)

$$M_H|_{\Lambda=1\,\text{TeV}} \gtrsim 50.8\,\text{GeV} + 0.64(m_t - 173.1\,\text{GeV}), \tag{37.}$$

which has already been surpassed by searches at LEP. For $\Lambda = M_{\text{Planck}}$, the lower bound rises to

$$M_H|_{\Lambda=M_{\text{Planck}}} \gtrsim 134\,\text{GeV}. \tag{38.}$$

Only noninteracting, or trivial, scalar field theories make sense on all energy scales. With restrictions, such theories can make sense up to a specified scale Λ, at which new physics comes into play. By analyzing the Q^2 evolution of the running quartic coupling in Equation 13, it is possible to

[2]The substantial literature on this topic may be traced from the state-of-the-art papers cited here.

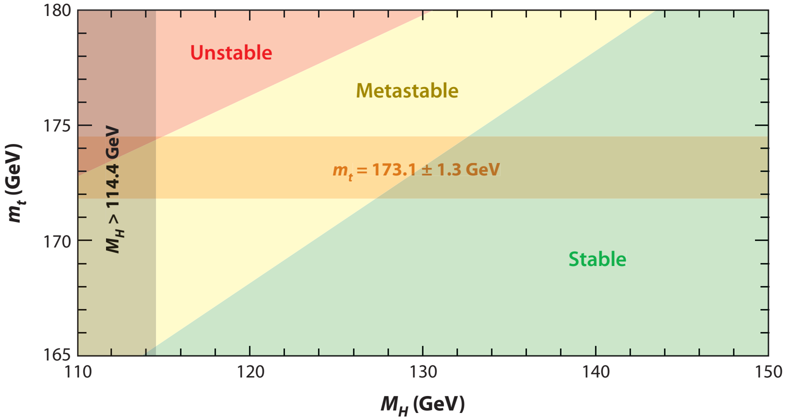


Figure 10

Metastability region of the Standard-Model vacuum in the (M_H, m_t) plane (151). The gray region at left indicates the LEP lower bound, $M_H > 114.4$ GeV. The horizontal orange band shows the measured top quark mass, $m_t = (173.1 \pm 1.3)$ GeV (101).

establish an upper bound on the coupling, and hence on the Higgs boson mass, at some reasonable scale accessible to experiment. A two-loop analysis leads to the bounds (147)

$$M_H|_{\Lambda=M_{\mathrm{Planck}}} \lesssim 180\,\mathrm{GeV} \qquad 39.$$

and

$$M_H|_{\Lambda=1\,\mathrm{TeV}} \lesssim 700\,\mathrm{GeV}. \qquad 40.$$

The electroweak theory could in principle be self-consistent up to very high energies, provided that the Higgs boson mass lies in the interval $134\,\mathrm{GeV} \lesssim M_H \lesssim 180\,\mathrm{GeV}$. If M_H lies outside this band, new physics will intervene at energies below the Planck (or unification) scale.

It is of considerable interest to use the techniques of lattice field theory to explore nonperturbative aspects of Higgs physics. What has been learned so far can be traced from References 148 and 149.

An informative perspective on the lower bound given in Equation 38 can be gained by relaxing the requirement that the electroweak vacuum correspond to the absolute minimum of the Higgs potential. It is consistent with observations for the ground state of the electroweak theory to be a false (metastable) vacuum that has survived quantum fluctuations until now. The relevant constraint is then that the mean time to tunnel from our electroweak vacuum to a deeper vacuum exceeds the age of the Universe, approximately 13.7 Gyr (150).

Figure 10 shows the (M_H, m_t) regions in which the Standard-Model vacuum is stable, acceptably long lived, or too short lived, as inferred from a renormalization-group-improved one-loop calculation of the tunneling probability at zero temperature (151). Present constraints on (M_H, m_t) suggest that we do not live in the unstable vacuum that would mandate new physics below the Planck scale. At the lowest permissible Higgs boson mass, this conclusion holds only at 68% CL.

4.2. Experimental Constraints on the Higgs Boson

We have seen in our discussion of evidence for the virtual influence of the Higgs boson in Section 3.5 that global fits, made within the framework of the standard electroweak theory, favor a light Higgs boson and exhibit some tension with direct searches. The LEP experiments, which

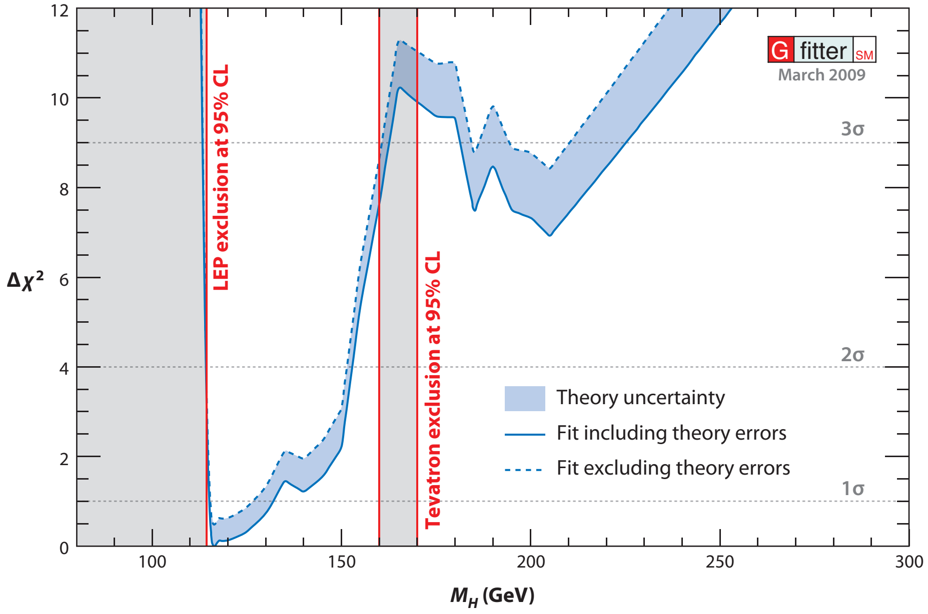


Figure 11

$\Delta\chi^2$ as a function of the Higgs boson mass for the Gfitter complete fit, taking account of direct searches at LEP and the Tevatron. The solid (dashed) blue line gives the results when including (ignoring) theoretical errors. The minimum $\Delta\chi^2$ of the fit, including theoretical errors, is used for both curves to obtain the offset-corrected $\Delta\chi^2$ (103).

focused on the $e^+e^- \rightarrow HZ^0$ channel, set a lower bound on the Standard-Model Higgs boson mass of $M_H > 114.4$ GeV at 95% CL (111, 152). The Tevatron experiments CDF and D0 also search for the Standard-Model Higgs boson, examining a variety of production channels and decay modes appropriate to different Higgs boson masses. The most recent combined result excluded the range 160 GeV $< M_H <$ 170 GeV at 95% CL (153, 112). See Reference 154 for an overview of past searches.

The disjoint exclusion regions from LEP and the Tevatron make it somewhat complicated to specify the remaining mass ranges favored for the standard-model Higgs boson. A useful example is shown in **Figure 11** (103). In the Gfitter analysis, at 2σ significance ($\approx$95% CL), the Standard-Model Higgs boson mass must lie in the interval 113.8 GeV $< M_H <$ 152.5 GeV.

The standard electroweak theory gives an excellent account of many pieces of data over a wide range of energies, and its main elements can be stated compactly. Nevertheless, it leaves too many gaps in our understanding for it to be considered a complete theory (cf. Section 5). We therefore have reason to consider extensions to the Standard Model, for which the Standard-Model fits to the electroweak measurements do not apply. Accordingly, healthy skepticism dictates that we regard the inferred constraints on M_H as a potential test of the Standard Model, not as rigid boundaries on where the agent of electroweak symmetry breaking must show itself.

Supersymmetric extensions of the electroweak theory entail considerable model dependency but yield high-quality fits to the precision data (155–157). On the one hand, bounds inferred from searches for the lightest *CP*-even Higgs boson h of the minimal supersymmetric Standard Model are somewhat less restrictive than for the standard-model Higgs boson. The tension between fits that prefer light masses and direct searches that disfavor a light Higgs boson is not present in the supersymmetric world. On the other hand, in its simplest form, the minimal supersymmetric Standard Model would be challenged if M_h exceeded approximately 135 GeV. A thorough discussion appears in Section 7.1 of Reference 158. A recent 25-parameter fit to the so-called phenomenological minimal supersymmetric Standard Model concludes that 117 GeV $\lesssim M_h \lesssim$ 129 GeV (157).

ATLAS: one of two general-purpose experiments for the LHC, located adjacent to CERN's main campus; see **http://atlasexperiment.org**

Compact Muon Solenoid (CMS): one of two general-purpose experiments for the LHC, located in Cessy, France. See **http://cms.cern.ch**

If new strong dynamics—rather than a perturbatively coupled elementary scalar—hides the electroweak symmetry, then the mass of the composite stand-in for the Higgs boson can range up to several hundred gigaelectronvolts. The same is true for Standard-Model fits that allow an extra generation of quarks and leptons (159, 160).

It would be prudent to search for the agent of electroweak symmetry breaking over the entire mass range allowed by general arguments, and this is what the LHC experiments will do. As an illustration, we next consider some elements of a broad search for the Standard-Model Higgs boson. This is a point of departure for more exotic searches. The search for the Higgs boson is now the province of the proton accelerators. The 2-TeV proton-antiproton Tevatron Collider is operating now, its integrated luminosity having surpassed 6 fb^{-1}, and the 14-TeV LHC at CERN will provide high-luminosity proton-proton collisions beginning in 2009.

4.3. Search for the Standard-Model Higgs Boson

The search for the Higgs boson has been a principal goal of particle physics for many years, so theorists and experimentalists have explored search strategies in great detail. The techniques in use at the Tevatron may traced from Reference 112, and the protocols foreseen for experiments at the LHC are detailed in the ATLAS (161) and CMS (162) performance documents.

Because the Standard-Model Higgs boson gives mass to the fermions and weak gauge bosons, it decays preferentially into the most massive states that are kinematically accessible. $H \to f\bar{f}$ decays, where fermion f occurs in N_c colors, proceed at a rate

$$\Gamma(H \to f\bar{f}) = \frac{G_F m_f^2 M_H}{4\pi\sqrt{2}} \cdot N_c \cdot \left(1 - \frac{4m_f^2}{M_H^2}\right)^{3/2}, \quad 41.$$

which is proportional to $N_c m_f^2 M_H$ as the Higgs boson mass becomes large. The partial width for decay into a W^+W^- pair is

$$\Gamma(H \to W^+W^-) = \frac{G_F M_H^2}{32\pi\sqrt{2}}(1-x)^{1/2}(4 - 4x + 3x^2), \quad 42.$$

where $x \equiv 4M_W^2/M_H^2$. Similarly, the partial width for decay into a pair of Z^0 bosons is

$$\Gamma(H \to Z^0 Z^0) = \frac{G_F M_H^2}{64\pi\sqrt{2}}(1-x')^{1/2}(4 - 4x' + 3x'^2), \quad 43.$$

where $x' \equiv 4M_Z^2/M_H^2$. The rates for decays into weak boson pairs are asymptotically proportional to M_H^3 and $\frac{1}{2}M_H^3$, respectively. In the final factors of Equations 42 and 43, $2x^2$ and $2x'^2$, respectively, arise from decays into transversely polarized gauge bosons. The dominant decays for large M_H are into pairs of longitudinally polarized weak bosons.

Branching fractions for decay modes that may hold promise for the detection of a Higgs boson are displayed in **Figure 12**. In addition to the $f\bar{f}$ and VV modes that arise at tree level, the plot includes the $\gamma\gamma$, $Z\gamma$, and two-gluon modes that proceed through loop diagrams.

The Higgs boson total width is plotted as a function of M_H in **Figure 13**. Below the W pair threshold, the Standard-Model Higgs boson is rather narrow, with $\Gamma(H \to \text{all}) \lesssim 1$ GeV. Far above the threshold for decay into gauge-boson pairs, the total width is proportional to M_H^3. As its mass increases toward 1 TeV, the Higgs boson becomes highly unstable, with a perturbative

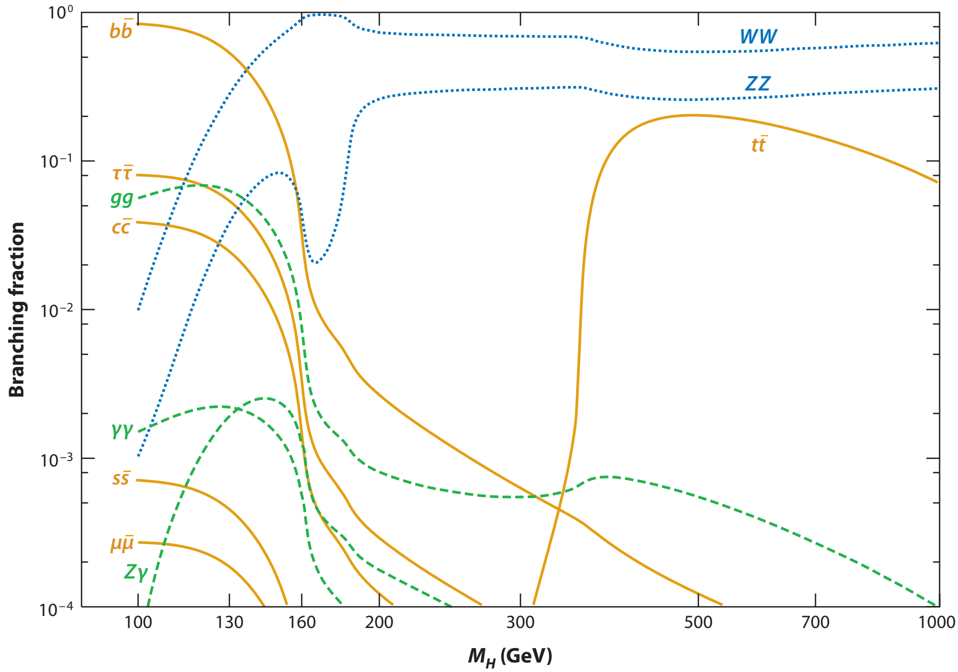


Figure 12

Branching fractions for prominent decay modes of the Standard-Model Higgs boson. Reproduced from Reference 163.

width approaching its mass. It would therefore be observed as an enhanced rate, rather than as a distinct resonance.

Cross sections for the principal reactions to be studied at the LHC are shown in **Figure 14**. The largest cross section for Higgs production at both the LHC and the Tevatron occurs in the reaction $p^{\pm}p \to H +$ anything, which proceeds by gluon fusion through heavy quark loops. (The shoulder in that cross section near $M_H = 400$ GeV reflects the behavior of the top quark loop.) A fourth generation of heavy quarks would raise the $gg \to H$ rate significantly, increasing the sensitivity of searches at the Tevatron and LHC.

For small Higgs boson masses, the dominant decay is into $b\bar{b}$ pairs, but the reaction $p^{\pm}p \to H +$ anything followed by the decay $H \to b\bar{b}$ is swamped by QCD production of $b\bar{b}$ pairs. Consequently, experiments must rely on rare decay modes (e.g., $\tau^+\tau^-$ or $\gamma\gamma$) with lower backgrounds or resort to different production mechanisms for which specific reaction topologies reduce backgrounds. Accordingly, the production of Higgs bosons in association with electroweak gauge bosons is receiving close scrutiny at the Tevatron. The rare $\gamma\gamma$ channel is seen as an important target for LHC experiments, if the Higgs boson is light. Fine resolution of the electromagnetic calorimeters is a prerequisite to overcoming Standard-Model backgrounds. At higher masses, the Tevatron experiments have exploited good sensitivity to the $gg \to H \to W^+W^-$ reaction chain to set their exclusion limits (112).

At the LHC, the multipurpose CMS and ATLAS detectors will make a comprehensive exploration of the Fermi scale, with high sensitivity to the Standard-Model Higgs boson reaching to 1 TeV. Current projections suggest that a few tens of femtobarns will suffice for a robust discovery (161, 162).

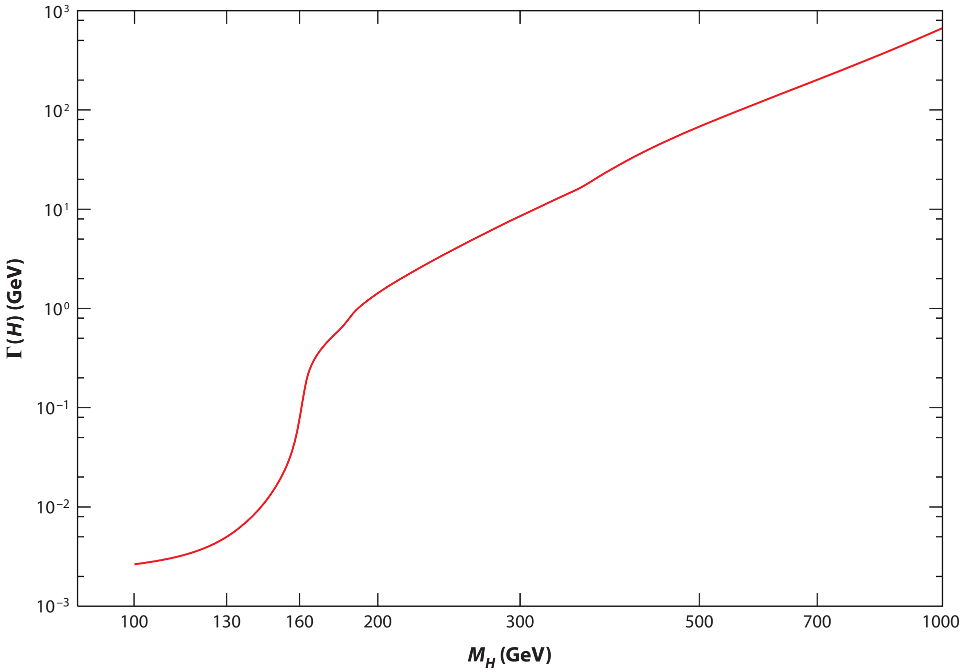


Figure 13

Total width of the Standard-Model Higgs boson versus mass. Reproduced from Reference 163.

Once the Higgs boson is found, it will be of great interest to map its decay pattern in order to characterize the mechanism of electroweak symmetry breaking. It is by no means guaranteed that the same agent hides electroweak symmetry and generates fermion mass. In Section 4.4.1, we recall that chiral symmetry breaking in QCD could hide the electroweak symmetry without generating fermion masses. Indeed, many extensions to the Standard Model significantly alter the decay pattern of the Higgs boson. In supersymmetric models, five Higgs bosons are expected, and the branching fractions of the lightest one may be very different from those presented in **Figure 12** (165).

Precise determinations of Higgs boson couplings is one of the strengths of the projected International Linear Collider (ILC) (166, 167), but the LHC will supply crucial clues to the origin of fermion masses. For example, a Higgs boson discovery in gluon fusion ($gg \to H$), signaled by the large production rate, would argue for a nonzero coupling of the Higgs boson to top quarks—an important qualitative conclusion. In time, and through comparison with other production and decay channels, it should be possible to constrain the $Ht\bar{t}$ coupling. With the LHC's large data sets, it is plausible that Higgs boson couplings could eventually be measured at levels that test the Standard Model and provide interesting constraints on extensions to the electroweak theory.

International Linear Collider (ILC): a project to design and construct an electron-positron collider initially reaching c.m. energy of 500 GeV; see **http://linearcollider.org**

4.4. Alternatives to the Higgs Mechanism

The Higgs mechanism shows how electroweak symmetry could be broken to reproduce the low-energy features of the real world, but it is not the only possibility. Before looking at some of the alternatives, let us recall that QCD hides the electroweak symmetry, although not in a way

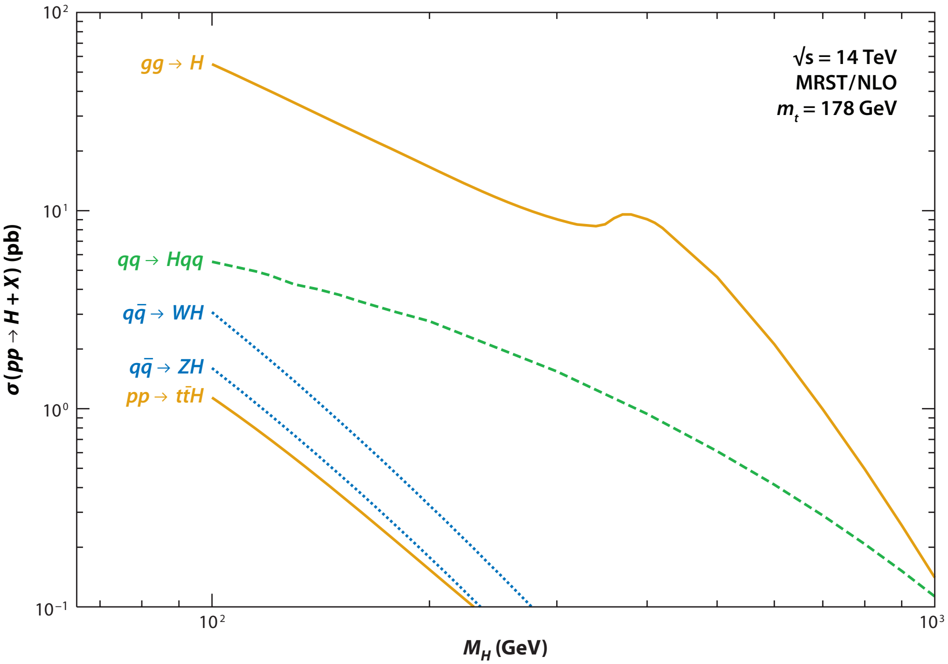


Figure 14

Higgs boson production cross sections in pp collisions at $\sqrt{s} = 14$ TeV, computed at next-to-leading order (NLO) using the MRST parton distributions (164). Reproduced from Reference 163.

that agrees with observation. A brief aside illuminates the importance of electroweak symmetry breaking by considering a Gedanken world without a specific mechanism to hide the electroweak symmetry. Then, I describe dynamical symmetry-breaking mechanisms modeled on the QCD example as well as new hypotheses that entail new global symmetries, new degrees of freedom, or supplementary dimensions of spacetime.

4.4.1. How quantum chromodynamics would hide electroweak symmetry. An analogy between electroweak symmetry breaking and the superconducting phase transition led to the insight of the Higgs mechanism. The macroscopic order parameter of the Ginzburg–Landau phenomenology, which corresponds to the wave function of superconducting charge carriers, acquires a nonzero vacuum expectation value in the superconducting state. Within a superconductor, the photon acquires a mass $M_\gamma = \hbar/\lambda_{\mathrm{L}}$, where the London penetration depth, λ_{L}, characterizes the exclusion of magnetic flux by the Meissner effect. In the particle-physics counterpart, auxiliary scalars introduced to hide the electroweak symmetry pick up a nonzero vacuum expectation value that gives rise to masses for the $W^{\pm}$ and Z^0.

A deeper look at superconductivity reveals an example of a gauge-symmetry-breaking mechanism that does not rely on introducing an ad hoc order parameter. In the microscopic Bardeen–Cooper–Schrieffer (BCS) theory (168), the order parameter arises dynamically through the formation of correlated states of elementary fermions, the Cooper pairs of electrons. Part of the beauty of the BCS theory is that the only new ingredient required is insight. The elementary fermions (electrons) and gauge interactions (QED) needed to generate the correlated pairs are already present in the case of superconductivity. This suggests that the electroweak symmetry might also be broken dynamically, without the need to introduce scalar fields. Indeed, QCD can be the source of electroweak symmetry breaking.

Consider an $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ theory of massless up and down quarks. Because the strong interaction is strong and the electroweak interaction is feeble, we may treat the $SU(2)_L \otimes U(1)_Y$ interaction as a perturbation. For vanishing quark masses, QCD displays an exact $SU(2)_L \otimes SU(2)_R$ chiral symmetry. At an energy scale $\sim\Lambda_{QCD}$, the strong interactions become strong and quark condensates of the form

$$\langle \bar{q}q \rangle \equiv \langle \bar{u}u + \bar{d}d \rangle \quad 44.$$

appear. The chiral symmetry is spontaneously broken to the familiar flavor symmetry isospin,

$$SU(2)_L \otimes SU(2)_R \to SU(2)_V, \quad 45.$$

because the left-handed and right-handed quarks communicate through $\langle \bar{q}q \rangle = \langle \bar{q}_R q_L + \bar{q}_L q_R \rangle$. Three Goldstone bosons appear, one for each broken generator of the original chiral invariance. These bosons were identified by Nambu (169) as three massless pions.

The broken generators are three axial currents whose couplings to pions are measured by the pion decay constant $f_\pi \approx 92.4$ MeV (45), which is measured by the charged-pion lifetime. When we turn on the electroweak interaction, the electroweak gauge symmetry is broken because the left-handed and right-handed quarks, now coupled through the $\langle \bar{q}q \rangle$ condensate, transform differently under $SU(2)_L \otimes U(1)_Y$ gauge transformations. The electroweak bosons couple to the axial currents and acquire masses of order $\sim g f_\pi$. The mass-squared matrix,

$$\mathcal{M}^2 = \begin{pmatrix} g^2 & 0 & 0 & 0 \\ 0 & g^2 & 0 & 0 \\ 0 & 0 & g^2 & gg' \\ 0 & 0 & gg' & g'^2 \end{pmatrix} \frac{f_\pi^2}{4}, \quad 46.$$

where the rows and columns correspond to W_1, W_2, W_3, and $\mathcal{A}$, has the same structure as the mass-squared matrix for gauge bosons in the standard electroweak theory.

Diagonalizing this matrix (Equation 46), we find that the photon, which corresponds (as in the Standard Model) to the combination $A = (g\mathcal{A} + g' b_3)/\sqrt{g^2 + g'^2}$, emerges massless. Two charged gauge bosons, $W^\pm = (b_1 \mp i b_2)/\sqrt{2}$, acquire mass-squared $M_W^2 = g^2 f_\pi^2/4$, and a neutral gauge boson $Z = (-g'\mathcal{A} + g b_3)/\sqrt{g^2 + g'^2}$ obtains $M_Z^2 = (g^2 + g'^2) f_\pi^2/4$. The ratio,

$$M_Z^2/M_W^2 = (g^2 + g'^2)/g^2 = 1/\cos^2\theta_W, \quad 47.$$

where θ_W is the weak mixing parameter, reproduces the Standard-Model result. The would-be massless pions disappear from the physical spectrum, becoming the longitudinal components of the weak gauge bosons. Here the symmetry breaking is dynamical and automatic; it can be traced, through spontaneous chiral symmetry breaking and confinement, to the asymptotic freedom of QCD.

Electroweak symmetry breaking determined by preexisting dynamics stands in contrast to the standard electroweak theory, in which spontaneous symmetry breaking results from the ad hoc choice of $\mu^2 < 0$ for the coefficient of the quadratic term in the Higgs potential. Despite the structural similarity to the Standard Model, the chiral symmetry breaking of QCD does not yield a satisfactory theory of the weak interactions. The masses acquired by the intermediate bosons are 2500 times smaller than required for a successful low-energy phenomenology; the W boson mass is only $M_W \approx 30$ MeV (170) because its scale is set by f_π. Moreover, QCD does not give masses to the fermions: The up and down quark and the electron all remain massless. See Section 3.7 for remarks on the logical separation between electroweak symmetry breaking and fermion mass generation.

4.4.2. If no Higgs mechanism shaped the world. Having recalled that QCD induces the breaking of electroweak symmetry through the formation of $\langle \bar{q} q \rangle$ condensates, it is worth considering how different the world would have been without a Higgs mechanism or a substitute on the real-world electroweak scale (171). Eliminating the Higgs mechanism does not alter the strong interaction, so QCD would still confine colored objects into hadrons. In particular, the gross features of nucleons derived from QCD—such as nucleon masses—would be little changed if the up and down quark masses were set to zero.

In the real world, the small $m_d > m_u$ mass difference overcomes the electromagnetic mass shift that would render the proton heavier than the neutron and results in $M_n - M_p \approx 1.293$ MeV, so that the neutron is unstable to β decay and the proton is the lightest nucleus. This contribution is absent if the quarks are massless.

However, the fact that electroweak symmetry is broken on the QCD scale means that the strength of the weak interactions would be similar to the strength of the strong nuclear force. The analog of the Fermi constant, G_F, is enhanced by nearly seven orders of magnitude. This has many consequences, including the acceleration of β-decay rates and the amplification of weak-interaction mass shifts that tend to make the neutron outweigh the proton. Because the theory lacks a Higgs boson, scattering among weak bosons becomes strongly coupled on the hadronic scale, following the analysis we reviewed in Section 4.1.

Should the proton be stable, or should compound nuclei be produced and survive to late times in this alternate universe, the infinitesimal electron mass would compromise the integrity of matter. The Bohr radius of a would-be atom would be macroscopic (if not infinite), valence bonding would have no meaning, and stable structures would not form. In seeking the agent of electroweak symmetry breaking, we hope to learn why the everyday world is as we find it: why atoms, chemistry, and stable structures can exist.

4.4.3. Dynamical symmetry breaking. The observation that QCD dynamically breaks electroweak symmetry (but at too low a scale) inspired the invention of analogous no-Higgs theories in which dynamical symmetry breaking is accomplished by the formation of a condensate of new fermions subject either to QCD itself or to a new, asymptotically free, vectorial gauge interaction (often termed technicolor) that becomes strongly coupled at the teraelectronvolt scale.

Within QCD, hypothetical exotic (color **6**, **8**, **10**, ...) quarks would interact more strongly than the normal color triplets, so the chiral-symmetry breaking in exotic quark sectors would occur at much larger mass scales than the standard chiral-symmetry breaking we have just reviewed. If those mass scales were sufficiently high, exotic quark condensates could break $SU(2)_L \otimes U(1)_Y \to U(1)_{em}$ dynamically and yield phenomenologically viable $W^{\pm}$ and Z^0 masses (172). No exotic quarks have yet been detected either by direct observation or in the evolution of the strong coupling constant, α_s (173).

Technicolor theories posit new technifermions that are subject to a new technicolor interaction. The technifermion condensates that dynamically break the electroweak symmetry produce masses for the $W^{\pm}$ and Z^0 bosons. Choosing the scale on which the technicolor interaction becomes strong so that the technipion decay constant is given by $F_\pi^2 = G_F\sqrt{2}$ reproduces the gauge-boson masses of the standard electroweak theory. Technicolor shows how the generation of intermediate boson masses could arise without fundamental scalars or unnatural adjustments of parameters. By replacing the elementary Higgs boson with an object that is composite on the electroweak scale, it also offers an elegant solution to the naturalness problem of the Standard Model, presented in Section 5.2.

However, simple technicolor does not explain the origin of quark and lepton masses because no Yukawa couplings are generated between Higgs fields and quarks or leptons. Consequently, technicolor serves as a reminder that particle physics confronts two problems of mass: (*a*) explaining the masses of the gauge bosons, which demands an understanding of electroweak symmetry breaking, and (*b*) accounting for the quark and lepton masses, which requires not only an understanding of electroweak symmetry breaking but also a theory of the Yukawa couplings that set the scale of fermion masses in the Standard Model.

To endow the quarks and leptons with mass, it is necessary to embed technicolor in a larger extended technicolor framework (136–138) containing degrees of freedom that communicate the broken electroweak symmetry to the (technicolor-singlet) Standard-Model fermions. Specific implementations of these ideas face phenomenological challenges pertaining to FCNC, the large top-quark mass, and precision electroweak measurements, but the idea of dynamical symmetry breaking remains an important alternative to the standard elementary scalar. For reviews and a summary of recent developments, see References 73–75, and 174).

Other suggestive work in the area of dynamical symmetry breaking also builds on the metaphor of the BCS theory of superconductivity but attributes a special role to quarks of the third generation or beyond. A particularly rich strategy, based on the notion that a top quark condensate drives electroweak symmetry breaking, was initiated in References 175–177. The idea that condensation of a strongly coupled fourth generation of quarks could trigger electroweak symmetry breaking is a lively area of contemporary research (178).

4.4.4. Other mechanisms for electroweak symmetry breaking. Very informative surveys of new approaches to electroweak symmetry breaking are given in References 179 and 180. Much model building has occurred around the proposition that the Higgs boson is a pseudo-Nambu-Goldstone boson of a spontaneously broken approximate global symmetry, with the explicit breaking of this symmetry collective in nature; that is, more than one coupling at a time must be turned on for the symmetry to be broken. These so-called little Higgs theories feature weakly coupled new physics at the teraelectronvolt scale (181, 182). When supplemented with a new symmetry known as T-parity, under which new heavy particles are odd and Standard-Model particles are even, the little Higgs theories can survive precision electroweak constraints and proffer a dark matter candidate (183).

New ways of thinking about electroweak symmetry breaking arise when we contemplate the possibility that spacetime has more than the canonical four dimensions. Among the possibilities are models without a physical Higgs scalar, in which electroweak symmetry is hidden by boundary conditions (184–186). The unitarity violation (cf. Section 4.1) that would be present in a theory without a Higgs boson is softened—deferred to energy scales well above 1 TeV—by the exchange of Kaluza (187)–Klein (188) (KK) excitations (189) of Standard-Model particles such as the W. In this case, the KK recurrences constitute the new physics on the 1-TeV scale required by the general argument.

Suppose instead that the electroweak gauge theory is itself formulated in more than four dimensions. From our four-dimensional perspective, components of the gauge fields along the supplemental directions will be seen as scalar fields with respect to the conventional four-dimensional coordinates (190, 191).

It is even conceivable that the electroweak phase transition is an emergent phenomenon arising from strong dynamics among the weak gauge bosons (192). If we take the mass of the Higgs boson to very large values (beyond 1 TeV in the Lagrangian of the electroweak theory), the scattering among gauge bosons becomes strong, in the sense that $\pi\pi$ scattering becomes strong on the gigaelectronvolt scale, as discussed in Section 4.1. In that event, it is reasonable to speculate that resonances form among pairs of gauge bosons, that multiple production of gauge bosons becomes commonplace, and that resonant behavior could hold the key to understanding what hides the electroweak symmetry.

5. INCOMPLETENESS OF THE ELECTROWEAK THEORY

For all its successes, the electroweak theory leaves many questions unanswered. It does not explain the negative coefficient $\mu^2 < 0$ of the quadratic term in Equation 15 required to hide the electroweak symmetry, and it merely accommodates, but does not predict, fermion masses and mixings. The CKM framework describes what we know of *CP* violation but does not explain its origin. The discovery of neutrino flavor mixing, with its implication that neutrinos have mass, calls for an extension of the electroweak theory set out in Section 2. Moreover, an elementary Higgs sector is unstable against large radiative corrections. A pervasive nonzero vacuum expectation value for the Higgs field implies a uniform energy density of the vacuum that seems incompatible with observations. Neutrinos are the only dark matter candidates within the Standard Model. They appear to contribute only a small share of the inferred dark matter energy density as relativistic (hot) dark matter, not the cold dark matter required for structure formation in the early Universe. The *CP* violation observed in the quark sector, in accord with the CKM paradigm, seems far too small to account for the excess of matter over antimatter in the Universe.

5.1. The Problem of Identity

The structure of the $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ Standard Model of QCD and the electroweak theory, as described in Section 2, can be written down in a few lines, but to calculate physical processes within the Standard Model—to apply the Standard Model to the real world—we need to specify at least 26 parameters. As tokens for the coupling parameters of the three factors of the gauge group, we may choose the strong coupling constant α_s, the fine structure constant α_{em}, and the weak mixing parameter $\sin^2\theta_W$. Two parameters are required to specify the shape of the Higgs potential in Equation 15. Additionally, there are six quark masses and four parameters (three mixing angles and the *CP*-violating phase) of the CKM matrix (Equation 5). The charged leptons and (massive) neutrinos add six more mass parameters, three more mixing angles, and one more *CP*-violating phase. Adding the (QCD) vacuum phase implicated in the strong *CP* problem brings the total to 26. Two more *CP*-violating phases enter if the neutrinos are their own antiparticles (Majorana particles). At least 20 of these parameters are related to the physics of flavor.

The operational question—"What determines the masses and mixings of the quarks and leptons?"—can be restated more evocatively: "What makes a top quark a top quark, an electron an electron, and a neutrino a neutrino?" It is not enough to answer "the Higgs mechanism" because the fermion masses are a very enigmatic element of the electroweak theory. Once the electroweak symmetry is hidden, the electroweak theory accommodates fermion masses, but the values of the masses are set by the apparently arbitrary couplings of the Higgs boson to the fermions (cf. **Figure 9**). Nothing in the electroweak theory is ever going to prescribe those couplings. It is not that the calculation is technically challenging; it is that there is no calculation. Neutrino masses can be generated through Yukawa couplings and in new ways as well because the neutrino may be its own antiparticle (193).

Within the standard electroweak theory, it is not only the fermion masses, but also the mixing angles that parameterize the mismatch between flavor eigenstates and mass eigenstates, that are set by the Yukawa couplings. The family relationships captured in the (CKM) quark-mixing matrix are displayed in the ternary plot in the left pane of **Figure 15**. The coordinates are given by the

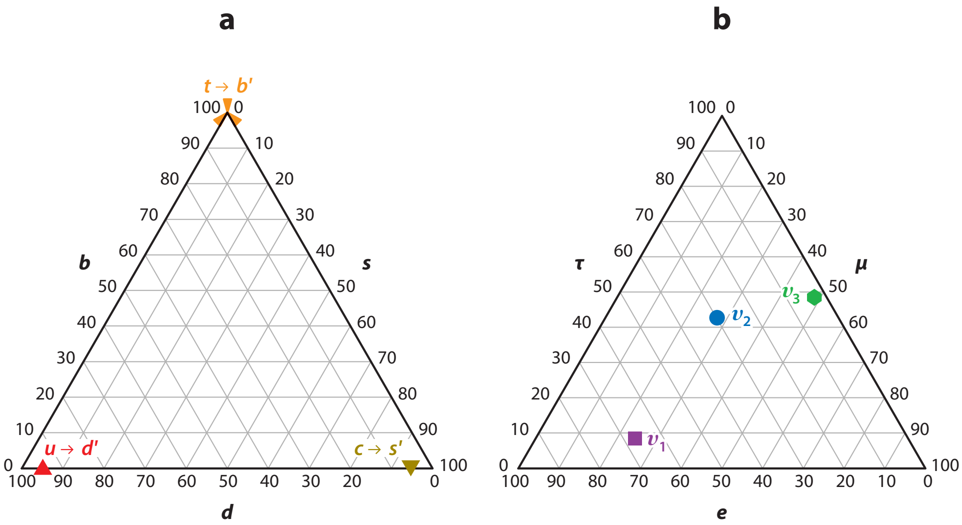


Figure 15

(*a*) *d*, *s*, and *b* composition of the quark flavor eigenstates d' (*red* Δ), s' (*brown* ∇), and b' (*orange tripod*). (*b*) ν_e, ν_μ, and ν_τ flavor content of the neutrino mass eigenstates ν_1, ν_2, and ν_3. The symbols denote central values, neglecting *CP* violation in the lepton sector, and with the "small" mixing angle taken as $\theta_{13} = 10°$ (194).

squares of the CKM matrix elements in each row of Equation 27. The *u* quark couples mostly to *d*, *c* mostly to *s*, and *t* almost exclusively to *b*.

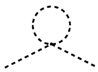
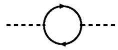
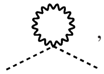
Our current knowledge of neutrino oscillations suggests the flavor content of the neutrino mass eigenstates depicted in **Figure 15*b***. The pattern is very different from that of the quark sector: The mass eigenstate ν_3 consists of nearly equal parts of ν_μ and ν_τ, perhaps with a trace of ν_e, whereas ν_2 contains similar amounts of ν_e, ν_μ, and ν_τ, and ν_1 is rich in ν_e, with approximately equal minority parts of ν_μ and ν_τ. Here ν_1 is the lighter of the solar pair, ν_2 is its heavier solar partner, and ν_3 lies either above (normal hierarchy) or below (inverted hierarchy) the solar pair in mass.

The exciting prospect is that quark and lepton masses, mixing angles, and subtle differences in the behavior of particles and their antiparticles put us in contact with physics beyond the Standard Model. One important step toward understanding will be to ascertain whether the Higgs boson is indeed the agent behind fermion mass. Another will be to determine whether the light neutrinos are in fact their own antiparticles, as would be signaled by the observation of neutrinoless double-beta decay. Perhaps we find it hard to decode the message in the fermion masses and mixings because we are only seeing part of a larger picture, and perhaps we must discover the spectrum of a new kind of matter—a fourth generation, or superpartners, or something entirely different—before it all begins to make sense.

5.2. The Problem of Widely Separated Scales

For all its triumphs, the standard electroweak theory has many shortcomings that lead physicists to surmise that the theory as we know it cannot be complete. Let us review some specific problems and suggested resolutions.

5.2.1. The hierarchy problem. Beyond the classical approximation, scalar mass parameters receive quantum corrections from loops that contain particles of spins $J = 0, \frac{1}{2}$, and 1, symbolically

$$M_H^2(p^2) = M_H^2(\Lambda^2) + $$  $$+$$  $$+$$  ,

where Λ defines a reference scale at which the value of M_H^2 is known. The dashed lines represent the Higgs boson, solid lines with arrows represent fermions and antifermions, and wavy lines represent gauge bosons. The quantum corrections that determine the running mass lead potentially to divergences,

$$M_H^2(p^2) = M_H^2(\Lambda^2) + \mathcal{C}g^2 \int_{p^2}^{\Lambda^2} dk^2 + \cdots, \tag{48.}$$

where g is the coupling constant of the theory, and the coefficient $\mathcal{C}$ is calculable in any particular theory. The loop integrals appear to be quadratically divergent, $\propto \Lambda^2$. In the absence of new physics, the reference scale Λ would naturally be large. If the fundamental interactions are described by QCD and the electroweak theory, then a natural reference scale is the Planck mass, $\Lambda \sim M_{\text{Planck}} = (\hbar c \,/\, G_{\text{Newton}})^{1/2} \approx 1.2 \times 10^{19}$ GeV. In a unified theory of the strong, weak, and electromagnetic interactions, a natural scale is the unification scale, $\Lambda \sim U \approx 10^{15} - 10^{16}$ GeV. Both estimates are very large compared to the electroweak scale, and so they imply a very long range of integration.

For the mass shifts induced by quantum corrections to remain modest, either something must limit the range of integration or new physics must damp the integrand. The challenge of preserving widely separated electroweak and reference scales in the presence of quantum corrections is known as the hierarchy problem. Unless we suppose that $M_H^2(\Lambda^2)$ and the quantum corrections are finely tuned to yield $M_H^2(p^2) \lesssim (1\ \text{TeV})^2$, some new physics—a new symmetry or new dynamics—must intervene at an energy of approximately 1 TeV to tame the integral in Equation 48.

Let us review the argument for the hierarchy problem: The unitarity argument (cf. Section 4.1) showed that new physics must be present on the 1-TeV scale in the form of either a Higgs boson or other new phenomena. However, a low-mass Higgs boson is imperiled by quantum corrections. New physics not far above the 1-TeV scale could bring the reference scale Λ low enough to mitigate the threat. That is what happens in models of large (195, 196) or warped (197, 198) extra dimensions (199), in which either (*a*) M_{Planck} is seen as a mirage, on the basis of a mistaken extrapolation of Newton's law of gravitation to very short distances, or (*b*) a new cutoff emerges, set by the scale of the extra dimension.

If the reference scale is indeed very large, then either various contributions to the Higgs boson mass must be precariously balanced or new physics must control the contribution of the integral in Equation 48. It is important to keep in mind that fine-tuning, perhaps guided by environmental selection, might be the way of the world (200). However, experience teaches us to be alert for symmetries or dynamics behind precise cancellations.

A new symmetry, not present in the Standard Model, could resolve the hierarchy problem. Exploiting the fact that fermion loops contribute with an overall minus sign relative to boson loops (because of Fermi statistics), supersymmetry (158, 165) balances the contributions of fermion and boson loops. In unbroken supersymmetry, the masses of bosons are degenerate with those of their fermion counterparts, so the cancellation is exact. If supersymmetry is present in our world, it must be broken. The contribution of the integrals may still be acceptably small if the fermion-boson mass splittings ΔM are not too large. The condition that $g^2 \Delta M^2$ be small enough leads to the requirement that superpartner masses be less than about 1 TeV. It is provocative to note that, with superpartners at $\mathcal{O}(1\ \text{TeV})$, the $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ coupling constants run to a common value at a unification scale of about 10^{16} GeV (201).

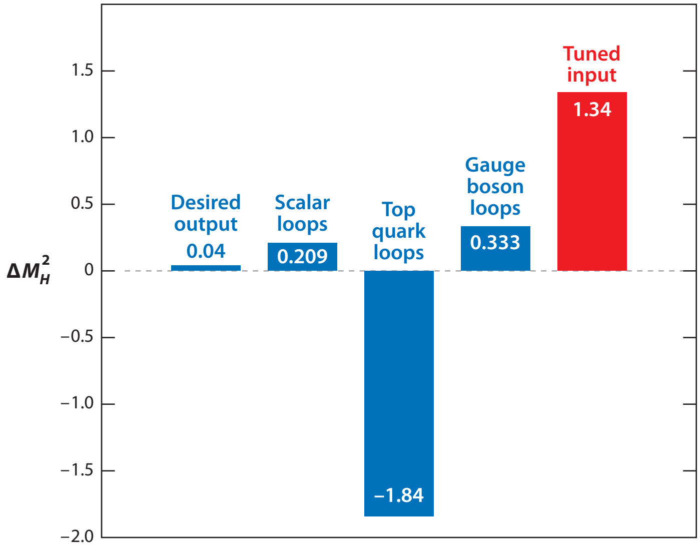


Figure 16

Relative contributions to ΔM_H^2 for a modest value of the cutoff parameter, $\Lambda = 5$ TeV, in Equation 48.

Theories of dynamical symmetry breaking (cf. Section 4.4.3) offer a second solution to the problem of the enormous range of integration in Equation 48. In technicolor models, the Higgs boson is composite, and its internal structure comes into play on the scale of its binding, $\Lambda_{\mathrm{TC}} \simeq \mathcal{O}(1\ \mathrm{TeV})$. The integrand is damped, the effective range of integration is cut off, and mass shifts are kept under control.

A recurring hope among theorists has been the notion that the Higgs boson might be naturally light because it is the pseudo-Nambu-Goldstone boson of some approximate global symmetry. Little Higgs models (181–183) introduce additional gauge bosons, vector-like quarks, and scalars on the teraelectronvolt scale. These conspire, thanks to a global symmetry, to cancel the quadratic divergences in Equation 48 that result from loops of Standard-Model particles and defer the hierarchy problem to approximately 10 TeV. In contrast to supersymmetry, the cancellations arise from loops containing particles of the same spin. In so-called twin Higgs models (202), the new states do not carry Standard-Model charges. The new physics at ~10 TeV raises impediments to conventional hopes for perturbative unification of the strong, weak, and electromagnetic interactions.

5.2.2. Tension between global fits and no new phenomena. A fine-tuning problem may be seen to arise even when the scale Λ is not extremely large. What has been termed the LEP paradox (203, 204) refers to a tension within the precise measurements of electroweak observables carried out at LEP and elsewhere. On the one hand, the global fits summarized in **Figure 6** point to a light Standard-Model Higgs boson. On the other hand, a straightforward effective-operator analysis of possible beyond-the-Standard-Model contributions to the same observables gives no hint of any new physics—of the kind needed to resolve the hierarchy problem—below approximately 5 TeV.

Figure 16 shows that, even with a cutoff $\Lambda = 5$ TeV, a careful balancing act is required to maintain a small Higgs boson mass in the face of quantum corrections, within the Standard Model, for which

$$\delta M_H^2 = \frac{G_{\mathrm{F}}\Lambda^2}{4\pi^2\sqrt{2}}(6M_W^2 + 3M_Z^2 + M_H^2 - 12m_t^2). \quad 49.$$

The chief cause for concern is the large contribution from the top quark loop,

$$\delta M_H^2\big|_{t-\text{loop}} \approx -\frac{3G_F}{\sqrt{2}\pi^2} m_t^2 \Lambda^2 \approx -0.075\,\Lambda^2. \qquad 50.$$

We are left to ask what enforces the balance, or how we might be misreading the evidence.

5.3. The Vacuum Energy Problem

The cosmological-constant problem—why empty space is so nearly massless—is one of the great mysteries of science (205, 206). It is the reason that gravity has weighed on the minds of electroweak theorists, despite the utterly negligible role that gravity plays in particle reactions. Recall that the gravitational attraction between an electron and proton is 41 orders of magnitude smaller than the electrostatic attraction at the same separation.

At the vacuum expectation value $\langle\phi\rangle_0$ of the Higgs field, the (position-independent) value of the Higgs potential is

$$V(\langle\phi^\dagger\phi\rangle_0) = \frac{\mu^2 v^2}{4} = -\frac{|\lambda| v^4}{4} < 0. \qquad 51.$$

Identifying $M_H^2 = -2\mu^2$, we see that the Higgs potential contributes a uniform vacuum energy density,

$$\varrho_H \equiv \frac{M_H^2 v^2}{8}. \qquad 52.$$

From the perspective of general relativity, this amounts to adding a cosmological constant, $\Lambda = (8\pi G_N/c^4)\varrho_H$, to Einstein's equation, where G_N is Newton's gravitational constant (207–209).

Recent observations of the accelerating expansion of the Universe (210, 211) raise the intriguing possibility that the cosmological constant may be different from zero, but the essential fact is that the observed vacuum energy density must be very small indeed (150),

$$\varrho_{\text{vac}} \lesssim 10^{-46}\,\text{GeV}^4 \approx (\text{a few millielectronvolts})^4. \qquad 53.$$

Therein lies the puzzle: If we take $v = (G_F\sqrt{2})^{-\frac{1}{2}} \approx 246\,\text{GeV}$ and insert the current experimental lower bound (111) $M_H \gtrsim 114.4\,\text{GeV}$ into Equation 52, we find that the Higgs field's contribution to the vacuum energy density is

$$\varrho_H \gtrsim 10^8\,\text{GeV}^4, \qquad 54.$$

some 54 orders of magnitude larger than the upper bound inferred from the cosmological constant. This mismatch has been a source of dull headaches for more than three decades.

The problem is still more serious in a unified theory of the strong, weak, and electromagnetic interactions, in which other (heavy!) Higgs fields have nonzero vacuum expectation values that may give rise to still larger vacuum energies. At a fundamental level, we can therefore conclude that a spontaneously broken gauge theory of the strong, weak, and electromagnetic interactions—or merely of the electroweak interactions—cannot be complete. The vacuum energy problem must be an important clue—but to what?

The tentative evidence for a nonzero cosmological constant recasts the problem in two important ways. First, instead of looking for a principle that would forbid a cosmological constant, perhaps a symmetry principle that would set it exactly to zero, we may be called upon to explain a tiny cosmological constant. Second, if the interpretation of the accelerating expansion in terms of dark energy is correct, we now have observational access to some new stuff whose equation of state and other properties we can try to measure. Maybe that will give us the clues that we need to solve this old problem and to understand how it relates to the electroweak theory.

5.4. Lacunae

The electroweak theory is unresponsive to some questions that are inspired by observations of the Universe at large.

5.4.1. Dark matter. The rotation curves of spiral galaxies and supporting evidence from the cosmic microwave background and large-scale structure point to dark matter that makes up 25% of the Universe's energy density (150). An appealing interpretation is that the dark matter is composed of one or more neutral relics from the early Universe. Within the Standard Model, the only candidates are neutrinos, for which the weight of experimental and observational evidences argues for masses smaller than approximately 1 eV.

Using the calculated number density of 56 cm^{-3} for each ν and $\bar{\nu}$ flavor in the current Universe, we can deduce the neutrino contribution to the mass density, expressed in units of the critical density, as $\rho_c \equiv 3H_0^2/8\pi G_\text{N} = 1.05h^2 \times 10^4\ \text{eV cm}^{-3} = 5.6 \times 10^3\ \text{eV cm}^{-3}$, where H_0 is the Hubble parameter now, G_N is Newton's constant, and I have taken the reduced Hubble constant to be $h = 0.73$ (212). Neutrinos contribute a normalized energy density $\Omega_\nu \gtrsim (1.2, 2.2) \times 10^{-3}$ for the (normal, inverted) spectrum, and no more than 10% of critical density, if the lightest neutrino mass approaches 1 eV. Neutrinos are not, however, candidates for the cold dark matter (nonrelativistic at the time of structure formation) that is favored by scenarios for structure formation in the Universe (213–215).

5.4.2. Baryon asymmetry of the Universe. Why does matter dominate over antimatter in the observable Universe (216)? Observations indicate that the density of antibaryons is negligible, whereas the average density of baryonic matter is characterized by the baryon-to-photon ratio,

$$\eta \equiv \frac{n_B - n_{\bar{B}}}{n_\gamma} = (6.14 \pm 0.25) \times 10^{-10}, \quad 55.$$

where n_B, $n_{\bar{B}}$, and n_γ are respectively the baryon, antibaryon, and photon number densities. Cosmological observations, anchored by the WMAP measurements of the Doppler peaks of temperature fluctuations in the cosmic microwave background radiation (150), imply that the current normalized baryon energy density is

$$\Omega_B = 0.0456 \pm 0.0015, \quad 56.$$

from which one can infer

$$\eta = (6.22 \pm 0.19) \times 10^{-10}. \quad 57.$$

Why is the ratio not zero? In an inflationary cosmology, conventional processes should produce equal numbers of baryons and antibaryons.

Three conditions are required to generate a baryon asymmetry out of neutral initial conditions (217): (*a*) the existence of fundamental processes that violate baryon number, (*b*) microscopic *CP* violation, and (*c*) departure from thermal equilibrium during the epoch in which baryon number–violating processes were important. A clear and compact survey of our current understanding of the baryon number of the Universe appears in Reference 218. How well does the electroweak theory respond?

The nonequilibrium condition is met by the expanding Universe. The electroweak theory does contain *CP* violation, in the CKM framework. At the level of perturbation theory, the electroweak theory conserves baryon number B and lepton number L, but that is not the case in the nonperturbative realm. Weak $SU(2)_\text{L}$ instantons violate B and L, conserving $B - L$, but have a negligibly small effect at temperatures T much lower than the electroweak scale $v \approx 246$ GeV. Their contributions to physical processes are suppressed by the factor $\exp(-8\pi^2/g^2)$ at zero temperature, where g is the $SU(2)_L$ gauge coupling (219). For $T \gtrsim v$, the sphalerons tend to erase a preexisting baryon asymmetry of the Universe and could under some conditions generate a significant baryon asymmetry. Our best assessment is that electroweak baryogenesis, within the Standard Model, falls well short of explaining the ratio in Equation 57. Some new physics, beyond the Standard Model, is required. A popular hypothesis is leptogenesis, in which the baryon asymmetry of the Universe is produced from a lepton asymmetry generated in the decays of a heavy sterile neutrino (220, 221).

5.4.3. Quantization of electric charge. The proton and electron charges balance to an astonishing degree (45):

$$|Q_p + Q_e| < 10^{-21}|Q_e|. \qquad 58.$$

If there were no connection between quarks and leptons, given that quarks make up the proton, then the balance of the proton and electron charge would simply be a remarkable coincidence, which seems an unsatisfying explanation. Some principle must relate the charges of the quarks and the leptons. What is it? An appealing strategy is to assign quarks and leptons to extended families. This is the approach of unified theories of the strong, weak, and electromagnetic interactions. It carries the implication of interactions that transform quarks into leptons, which has consequences for proton decay and baryogenesis. The idea that QCD and the electroweak theory might be unified is made plausible by the fact that both are gauge theories with similar mathematical structures.

Another encouragement to consider quark-lepton unification comes from the electroweak theory itself. A world governed by $SU(2)_L \otimes U(1)_Y$ interactions and composed only of quarks, or only of leptons, would be anomalous (13), in the technical sense that quantum corrections would break the gauge symmetry on which the theory is based. In our left-handed world, an anomaly-free electroweak theory is possible only if weak-isospin pairs of color-triplet quarks accompany weak-isospin pairs of color-singlet leptons. For these reasons, it is nearly irresistible to consider a unified theory that puts quarks and leptons into a single extended family.

5.4.4. Absence of gravity. The gravitational force is famously negligible in the realm of particle physics. For example, the gravitational attraction between a proton and electron is some forty-one orders of magnitude weaker than the Coulomb attraction, at the same separation. In the language of Feynman rules, dimensional analysis shows that the emission of a graviton is suppressed by a factor

$$E^\star / M_{\text{Planck}}, \qquad 59.$$

where $E^\star$ is a characteristic energy of the transition. The Planck mass ($M_{\text{Planck}} \equiv (\hbar c / G_{\text{Newton}})^{1/2} \approx 1.22 \times 10^{19}$ GeV) is a big number because Newton's constant is small in the appropriate units. Except in special circumstances, such as the excitation of many extradimensional modes (222), graviton emission need not be taken into account in particle physics.

However, we have seen in our discussion of the vacuum energy problem (cf. Section 5.3) that the relationship of the electroweak theory to gravitation cannot be ignored. The hierarchy problem (cf. Section 5.2) is acute, in part because the Planck scale is so distant from the electroweak scale. These connections, as well as the desire to understand why gravity is so much weaker than the $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ interactions, motivate efforts to integrate gravity with the Standard Model. When imagining how this could be achieved, it is important to bear in mind that we have probed the electroweak theory and QCD up to approximately 1 TeV, but we have tested the inverse-square law of gravity only down to distances shorter than 0.1 mm, corresponding to energies of 10 millielectronvolts!

6. THE NEW ERA

The Tevatron is expected to operate through 2011 and to produce a total of 10 fb^{-1} for analysis by the CDF and D0 collaborations. The experimenters are optimistic that a sample of that size will be sufficient—in the absence of a signal—to set a 95% exclusion limit for the Standard-Model Higgs boson over the entire range currently favored by the global fits discussed in Section 4.2. Barring a breakthrough in analysis techniques, discovery of the Standard-Model Higgs boson at 5σ significance is extremely unlikely at the Tevatron, unless the production rate should be enhanced (for example, by a fourth generation of quarks). At the interesting level of 3σ evidence, the situation is more promising. The experiments have quoted the odds of establishing "evidence" at approximately one in three for $120\,\mathrm{GeV} \lesssim M_H \lesssim 145\,\mathrm{GeV}$ and better than one in two for $M_H \lesssim 116\,\mathrm{GeV}$ and $150\,\mathrm{GeV} \lesssim M_H \lesssim 177\,\mathrm{GeV}$ (223; J. Konigsberg, private communication). At a minimum, we will know more about where the (Standard-Model) Higgs boson is not by the time the LHC Higgs search begins in earnest.

The parameters of the LHC are shaped by the imperative to make a thorough exploration of the 1-TeV scale, and so to elucidate the mechanism of electroweak symmetry breaking. However, the LHC is a discovery machine, broadly understood, not limited to the search for the agent of electroweak symmetry breaking. See Reference 224 for a general survey, and the ATLAS (161) and CMS (162) physics reports for the detector capabilities, with many specific illustrations.

6.1. Electroweak Questions for Large Hadron Collider Experiments

Will the new physics that we anticipate on the 1-TeV scale be a Higgs boson, in some guise, or new strong dynamics? If it is a Higgs boson, will there be one or several? Will it—or they—turn out to be elementary or composite? Is the Higgs boson indeed light, as anticipated by the global fits to electroweak precision measurements? Does the Higgs boson—the H—give mass only to the electroweak gauge bosons, or does it also endow the fermions with mass? Proceeding step by step, does the H couple to fermions (a large $t\bar{t}H$ coupling might be inferred from its production rate)? Are the branching fractions for decays into fermion pairs in accord with the Standard Model? A difficult follow-up question, should we find that the Higgs boson *is* responsible for fermion mass, is: What determines the masses and mixings of the fermions?

The Higgs bosons could couple to particles beyond those known in the Standard Model. Does the pattern of Higgs boson decays imply new physics? Will unexpected or rare decays of the Higgs boson reveal new kinds of matter? If more than one, apparently elementary, Higgs boson is found, will that be a sign for a supersymmetric generalization of the Standard Model, or for a different sort of two-Higgs-doublet model? What stabilizes the Higgs boson mass below 1 TeV? How can a light Higgs boson coexist with the absence of signals for new phenomena? Is electroweak symmetry breaking an emergent phenomenon connected with strong dynamics? Is electroweak symmetry breaking related to gravity through extra spacetime dimensions?

If the new physics observed on the teraelectronvolt scale is suggestive of new strong dynamics, how can we diagnose the nature of the new dynamics? What takes the place of a Higgs boson?

6.2. More New Physics on the Teraelectronvolt Scale?

The partial-wave unitarity argument reviewed in Section 4.1 indicates that a thorough exploration of the teraelectronvolt scale will produce important insights into the nature of electroweak symmetry breaking. At a strongly suggestive level, we have reason to expect additional new phenomena in this energy range.

The large gap between the electroweak scale and the unification scale or the Planck scale menaces the Higgs boson mass with large quantum corrections that would lift it far above 1 TeV. If the required small mass of the Higgs boson is not stabilized by fine-tuning, then new physics is needed on the teraelectronvolt scale. Familiar examples are (*a*) supersymmetry, with a spectrum of superpartners beginning to appear on the teraelectronvolt scale, and (*b*) technicolor, with its own spectrum of new technipions and W^+W^- resonances. The common characteristic of solutions to the hierarchy problem is that they lead naturally to the expectation of new physics on the teraelectronvolt scale. See References 225 and 226 for insightful surveys of some new-physics signatures to be anticipated at the LHC.

The evidence that cold dark matter is a significant component of the energy-density budget of the Universe is an independent argument for new phenomena on the teraelectronvolt scale. As with the hierarchy problem, the implication is highly suggestive, but not inevitable: A few-hundred-gigaelectronvolt particle that couples with weak-interaction strength could supply the missing mass (227).

To summarize, we have three indications for dramatic new developments on the teraelectronvolt scale: the requirement for a Higgs boson or new dynamics and the strong suggestions of new phenomena to solve the hierarchy problem and of particle dark matter. With the great discovery reach of the LHC, many other possibilities are open, including new heavy fermions and new force particles.

6.3. How Knowledge Might Accumulate

How our understanding progresses in light of information from the LHC depends, of course, on what nature has in store for us and how our attention is attracted this way or that in light of discoveries (see sidebar, Initial Goals for the LHC). However, the extensive studies carried out in preparation for the ATLAS and CMS experiments, and informed by the Tevatron experience, allow us to anticipate the sensitivity required for various potential observations and discoveries. It is implicit that understanding of the detectors progresses as data are accumulated (228, 229).

With an integrated luminosity of only 10 pb^{-1}, the experiments will begin to characterize event structure in the new energy regime, to measure jet and dijet spectra, and to study J/ψ, Υ, and $W^\pm$ production. Within the first hundred days of stable running, at approximately

INITIAL GOALS FOR THE LHC

The LHC and its detectors are instruments of unprecedented size, complexity, and capability. As commissioning proceeds, the first milestone for the experiments will be to recognize the familiar Standard-Model particles in the new setting. From the very first collisions, the experiments will cast a broad net for novel phenomena that might be revealed in their vast new domain in energy and sensitivity. In particular, they will search for new forces of nature, signaled by gauge bosons decaying into lepton pairs or dijets. New strong interactions might reshape our view of electroweak symmetry breaking.

In time, run strategies will be shaped by the goal of determining what breaks the electroweak symmetry—in shorthand, searching for the Higgs boson. Once the Higgs boson is found, the world will change as many specific questions about the details of electroweak symmetry breaking become accessible. An important yes-no question is whether the Higgs boson gives mass to fermions as well as to gauge bosons.

Along the way, the experiments will seek evidence for new matter particles, including dark matter candidates, and for supersymmetry, hidden spacetime dimensions, and the totally unexpected.

50 pb^{-1}, $Z^0 \to \ell^+\ell^-$ comes into view, and $t\bar{t}$ pairs should be observed. The latter observation will be an important milestone because top quark events provide subtle training grounds for detector algorithms that will allow tops to be identified as physics objects and represent a background that must be mastered for many new-physics searches. When the data set reaches approximately 100 pb^{-1}, incisive measurements of Standard-Model parameters come into reach.

A few early discoveries are possible for data sets of a few hundred inverse picobarns at $\sqrt{s} =$ 14 TeV: a Z' representing a new force of nature, with mass up to approximately 1 TeV, and light ($\approx$500 GeV) squarks and gluinos that would give evidence of supersymmetry. At approximately 1 fb^{-1}, Standard-Model Higgs boson physics opens up, first with discovery sensitivity for the channel $H \to ZZ \to \mu^+\mu^-\mu^+\mu^-$ with $M_H \approx 180$ GeV. Establishing a Higgs boson signal at low mass is more demanding: At $M_H \approx 115$ GeV, an integrated luminosity $\gtrsim$5 fb^{-1} will be required. At that point, combined channels from the experiments will cover the full allowed range of Standard-Model Higgs boson masses.

When the LHC data set exceeds roughly 10 fb^{-1}, a spin-two dilepton resonance characteristic of extra spacetime dimensions will be visible up to 1 TeV. An order of magnitude more gives a discovery reach for leptoquarks up to 1.5 GeV and a sensitivity to a compositeness scale $\Lambda^* = 30$ TeV. Integrated luminosity of 300 fb^{-1}, representing approximately five years of LHC experimentation, should suffice for the observation of teraelectronvolt-scale W^+W^- resonances and should expand the discovery reach for squarks and gluinos up to 2.5 TeV.

Our understanding will advance not only by discovery of new phenomena, but also by improvement of precision. As one example, **Figure 17** projects improvements in the global-fit constraints on the mass of the Standard-Model Higgs boson in light of measurements to be carried out at the LHC, and measurements that might be made at the proposed ILC, a 500-GeV electron-positron

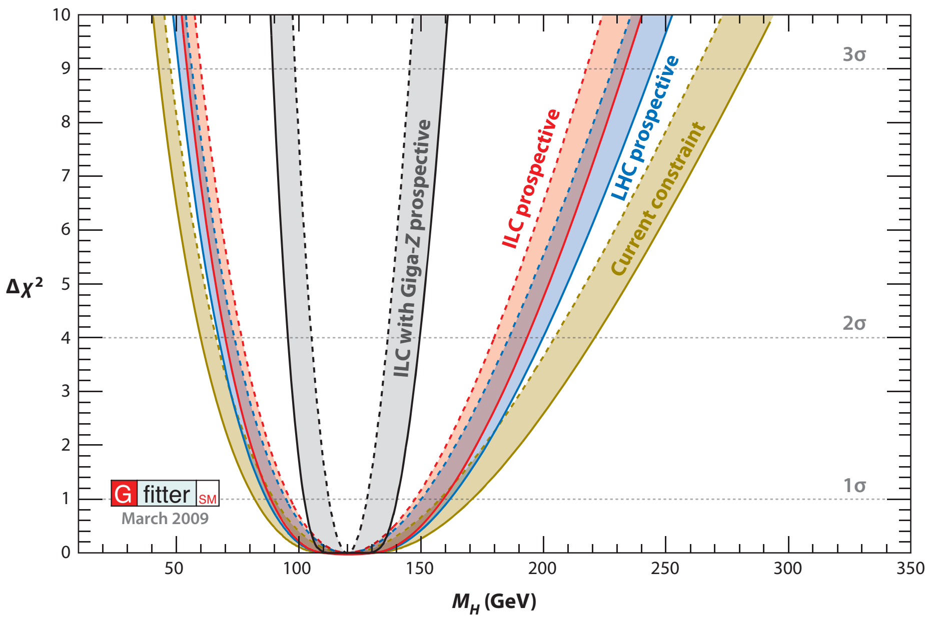


Figure 17

Gfitter constraints on the Higgs boson mass obtained for four future scenarios. Parabolas in $\Delta\chi^2$ are shown with their theoretical error bands. From wider to narrower: current constraint (*brown*), LHC expectation (*blue*), and ILC expectations excluding (*red*) and including (*gray*) the Giga-Z option (103).

collider. Increasing the sample of Z bosons by two orders of magnitude, the so-called Giga-Z option for a linear collider, further sharpens the constraint. (For surveys of the full ILC physics program, see References 166 and 167.) As was the case for the values of m_t inferred before the discovery of the top quark (cf. **Figure 4**), comparison of the global-fit constraints with the observed Higgs boson mass will be an incisive test of the Standard Model.

6.4. The Intensity Frontier

Historically, much of the motivation for the electroweak theory came from detailed measurements at low energies, and such experiments have led the validation of the CKM structure of the charged-current weak interaction and established the suppression of FCNC. The main imperative now is to explore the teraelectronvolt scale, to establish the mechanism for electroweak symmetry breaking. That task will soon pass from the Tevatron to the LHC. Many interesting questions remain in flavor physics. The importance of intensity-frontier experiments for reshaping our understanding of particle physics can be enhanced by their conversation with the LHC experiments that explore the teraelectronvolt scale.

New sensitivity can bring surprises. The FCNC examples we saw in Section 3.2 show that there are good opportunities for physics beyond the Standard Model to appear. Prospects for the study of B, D, and K decays are reviewed in Reference 230. The lepton sector is taking on greater interest following the discovery of neutrino mixing and the possibility that CP violation might be observable in neutrino interactions (231). Moreover, new searches for charged-lepton flavor violation and for CP violation in lepton dipole moments offer the possibility of dramatic discoveries (232). Interpretations of new phenomena observed at the LHC will be tested and refined through searches for the virtual effects of the new particles in rare processes studied at low energies. The LHC experiments, including LHC b, have their own role in flavor physics (233).

7. SUMMARY

Over the past two decades, experiments have tested the gauge sector and the flavor sector of the electroweak theory extensively, so that we may now regard the electroweak theory as a law of nature—subject, of course, to revision in the light of new evidence. Much of the experimental evidence was recounted in Section 3. The body of evidence is both broad and deep, and although it leaves room for physics beyond the Standard Model, it also constrains new physics in significant ways.

Experiments at the LHC will probe the electroweak symmetry–breaking sector on the 1-TeV scale, where we may also hope to find pointers to physics beyond the Standard Model. The clues in hand suggest that the agent of electroweak symmetry breaking represents a novel fundamental interaction operating on the Fermi scale. We do not know what that force is.

A leading possibility is that the agent of electroweak symmetry breaking is an elementary scalar, the Higgs boson of the electroweak Standard Model. Global fits to electroweak measurements indicate that the Standard-Model Higgs boson should be found with a mass not much more than 200 GeV. An essential step toward understanding the new force that shapes our world is, therefore, to search for the Higgs boson and to explore its properties by asking the questions highlighted in the Future Issues section.

We have seen in Section 3.5 that different sensitive observables prefer different values for the Higgs boson mass. We do not know whether that reflects an unexceptional scatter or whether it is a harbinger of physics beyond the Standard Model. In any case, it is important to search for the Higgs boson over the complete range of a priori acceptable values, and the ATLAS and CMS experiments will accomplish that.

As we detailed in Section 5, the Standard Model is incomplete. It shows how the masses of the quarks and leptons might arise, but it does not predict their values. It does not even give a qualitative understanding of why the quark-mixing parameters are small and hierarchical, nor why the pattern of neutrino mixing should be so different.

The hierarchy problem seems to require a solution in terms of dynamics or a symmetry—although fine-tuning is a logical possibility. The vacuum energy problem indicates that something essential is missing in our understanding. These are problems within the electroweak theory. Other shortcomings, including the absence of a dark matter candidate, speak to the limited reach of the electroweak theory.

We can be confident that the origin of gauge-boson masses will be understood on the teraelectronvolt scale. We do not know where we will decode the pattern of the Yukawa couplings that set the fermion masses. However, candidate solutions to the hierarchy problem entail new physics on the teraelectronvolt scale, and the WIMP (weakly interacting massive particle) solution to the dark matter question suggests a mass in the few-hundred-gigaelectronvolt range. These hints suggest that, in addition to the electroweak symmetry–breaking physics we confidently expect to see at the LHC, there is a great likelihood of more new phenomena.

The electroweak theory is a remarkable achievement. It gives a deeper understanding of two of the fundamental forces of nature—electromagnetism and the charged-current weak interaction—and adds the neutral-current weak interaction to the mix. It accounts for a wide variety of experimental measurements and has survived many tests as a quantum field theory. It meets the most important criteria for a good theory: We get more out than we put in, and it raises new and significant questions.

We are on the cusp of a new level of understanding, with the nature of electroweak symmetry breaking virtually certain to be revealed on the 1-TeV scale. At the same time, the incompleteness of the electroweak theory argues that we have much more to learn. Part of the high anticipation that attends the coming of the LHC is that the experimental opportunities on the 1-TeV scale involve distinct problems that could well be related, and indeed that may all be related through the electroweak theory. We will soon know how robust the connections are.

SUMMARY POINTS

1. The electroweak theory accounts for diverse phenomena over a prodigious range of distances or energies, but it leaves important questions unanswered. For example, it does not explain the masses of the quarks and leptons.
2. According to the Standard Model, the agent of electroweak symmetry breaking is an elementary scalar particle, the Higgs boson, but other mechanisms for electroweak symmetry breaking are possible. Experiments at the LHC will settle the issue.
3. Even if completed by the observation of a light Higgs boson, the electroweak theory raises questions. The hierarchy problem suggests more new physics on the teraelectronvolt scale. The vacuum energy problem requires an answer.
4. Tests of the electroweak theory also constrain the new physics that might be needed to extend the Standard Model. In particular, limits on FCNC interactions pose serious challenges.

5. Experiments at the LHC will be sensitive to a wide variety of new phenomena, including new forces that would imply extended symmetries, and new kinds of matter.
6. Proposed solutions to the hierarchy problem, the dark matter problem, and the nature of electroweak symmetry breaking suggest that these problems may be related.
7. Tentative interpretations of new particles and forces observed at the LHC must withstand the scrutiny of intensity-frontier experiments.
8. New insights into neutrino mixing and lepton-flavor violation may suggest new perspectives on quark mixings.

FUTURE ISSUES

1. What is the agent that hides the electroweak symmetry? Specifically, is there a Higgs boson? Might there be several?
2. Is the Higgs boson elementary or composite? How does the Higgs boson interact with itself? What triggers electroweak symmetry breaking?
3. Does the Higgs boson give mass to fermions, or only to the weak bosons? What sets the masses and mixings of the quarks and leptons?
4. What stabilizes the Higgs boson mass below 1 TeV?
5. Do the different behaviors of left-handed and right-handed fermions with respect to charged-current weak interactions reflect a fundamental asymmetry in the laws of nature?
6. What will be the next symmetry recognized in nature? Is nature supersymmetric? Is the electroweak theory part of some larger edifice?
7. Are there additional generations of quarks and leptons?
8. What resolves the vacuum energy problem?
9. Is electroweak symmetry breaking an emergent phenomenon connected with strong dynamics? Is electroweak symmetry breaking related to gravity through extra spacetime dimensions?
10. What lessons does electroweak symmetry breaking hold for unified theories of the strong, weak, and electromagnetic interactions?

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