Higgs Bosons, Electroweak Symmetry Breaking, and the Physics of the Large Hadron Collider

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The Large Hadron Collider, a $7\oplus7$ TeV proton-proton collider under construction at CERN (the European Laboratory for Particle Physics in Geneva), will take experiments squarely into a new energy domain where mysteries of the electroweak interaction will be unveiled. What marks the 1-TeV scale as an important target? Why is understanding how the electroweak symmetry is hidden important to our conception of the world around us? What expectations do we have for the agent that hides the electroweak symmetry? Why do particle physicists anticipate a great harvest of discoveries within reach of the LHC?

1 Introduction

Electromagnetism and the weak interactions share a common origin in the weak-isospin and weak-hypercharge symmetries described by the gauge group $\mathrm{SU(2)_L \otimes U(1)_Y}$, but their manifestations are very different. Electromagnetism is a force of infinite range, while the influence of the charged-current weak interaction responsible for radioactive beta decay only spans distances shorter than about 10^{-15} cm, less than 1% of the proton radius. We say that the electroweak gauge symmetry is spontaneously broken—hidden—to the $\mathrm{U(1)_{em}}$ phase symmetry of electromagnetism. How the electroweak gauge symmetry is hidden is one of the most urgent and challenging questions before particle physics.

The search for the agent that hides the electroweak symmetry is also one of the most fascinating episodes in the history of our quest to understand the material world. Over the next decade, experiments at the Large Hadron Collider (LHC) will lead us to a new understanding of questions that are both simple and profound: Why are there atoms? Why chemistry? What makes stable structures possible? Uncovering the answers to those questions may even bring new insight into "What makes possible the prerequisites for life?" A goal of this Article is to link these questions to the electroweak theory, and to the explorations soon to come at the LHC.

Within the standard electroweak theory, the agent of electroweak symmetry breaking is posited to be a single elementary scalar particle known as the Higgs boson, and so "the search for the Higgs boson" is a common token for the campaign to understand the origins of electroweak symmetry breaking. Such a shorthand is fine—so long as it is not taken to define a very limited menu of opportunities for discovery. As we embark upon the LHC adventure, we will need open and prepared minds!

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It is often repeated that the discovery of the Higgs boson will reveal the origin of all mass in the Universe. This statement is deeply incorrect—even if we restrict our attention to the luminous matter that is made of familiar stuff.[1] We can see why the familiar tagline is not right, even before we have reviewed precisely what we mean by the Higgs boson.

At each step down the quantum ladder, we understand mass in different terms. In quotidian experience, the mass of an object is the sum of the masses of its parts. At a level we now find so commonplace as to seem trivial, we understand the mass of any atom or molecule in terms of the masses of the atomic nuclei, the mass of the electron, and quantum electrodynamics.[1] And in precise and practical—if not quite "first-principle"—terms, we understand the masses of all the nuclides in terms of the proton mass, the neutron mass, and our knowledge of nuclear forces.

What about the proton and neutron masses? We have learned from Quantum Chromodynamics, the gauge theory of the strong interactions, that the dominant contribution to the light-hadron masses is not the masses of the quarks of which they are constituted, but the energy stored up in confining the quarks in a tiny volume (2). Indeed, the masses m_u and m_d of the up and down quarks are only a few MeV. The quark-mass contribution to the 939-MeV mass of an isoscalar nucleon (averaging proton and neutron properties) is only (3)

$$M_N^q = 3\,\frac{m_u + m_d}{2} = (7.5 \text{ to } 16.5) \text{ MeV} , \tag{1}$$

no more than 2%. Hadrons such as the proton and neutron thus represent *matter of a novel kind.* In contrast to macroscopic matter, atoms, molecules, and nuclei, the mass of a hadron is not equal to the sum of its constituent masses (up to small corrections for binding energy). The quark masses do account for an important detail of our world: The counterintuitive observation that the neutral neutron (udd) is more massive than the charged proton (uud) by 1.29 MeV is explained by the fact that m_d exceeds m_u by enough to overcome the proton's greater electromagnetic self-energy.

Our most useful tool in the strong-coupling regime is QCD formulated on a spacetime lattice. Calculating the light-hadron spectrum from first principles has been one of the main objectives of the lattice program, and important strides have been made recently. For example, the CP-PACS (Tsukuba) Collaboration's quenched calculation (no dynamical fermions) matches the observed light-hadron spectrum at the 10% level (4). Though small, the discrepancy is larger than the statistical and systematic uncertainties, and so is interpreted as a shortcoming of the quenched approximation. The unquenched results now emerging should improve the situation further (5), and give us new insights into how well—and why!—the simple quark model works.

The successful calculation of the hadron spectrum is a remarkable achievement for the theory of quantum chromodynamics and for lattice techniques. In identifying the energy of quark confinement as the origin of the nucleon mass, *we have explained nearly all the visible mass of the Universe,* since the luminous matter is essentially made of protons and neutrons in stars. To excellent approximation, that visible mass of the Universe arises from QCD—not from the Higgs boson.[2]

The Higgs boson and the mechanism of electroweak symmetry breaking are nevertheless of capital importance in shaping our world, accounting for the masses of the weak-interaction force particles and—at least in the standard electroweak theory—giving masses to the quarks and leptons. To develop that importance, we shall begin by sketching the electroweak theory and evoking its successes. Then we will address the key question: What would the world be like if there were no Higgs mechanism?

Once having established that the character of electroweak symmetry breaking is a compelling issue, we will consider where the crucial information should be found. The electroweak theory itself points to the

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[1]We do not know the nature of the dark matter in the Universe, so cannot yet explain how the mass of the dark-matter particles arises.
[1]In Dirac's 1929 formulation (1), "The underlying physical laws necessary for the mathematical theory of ...the whole of chemistry are thus completely known, and the difficulty is only that the exact application of these laws leads to equations much too complicated to be soluble."
[2]The standard model of particle physics (with its generalization to a unified theory of the stong, weak, and electromagnetic interactions) has taught us many fascinating interrelations, including the effect of heavy-quark masses on the low-energy value of the strong coupling constant, which sets the scale of the light-hadron masses. For a quick tour, see my *Physics Today* article on the top quark (6); Bob Cahn's *RMP* Colloquium (7) takes a more expansive look at connections within the standard model.
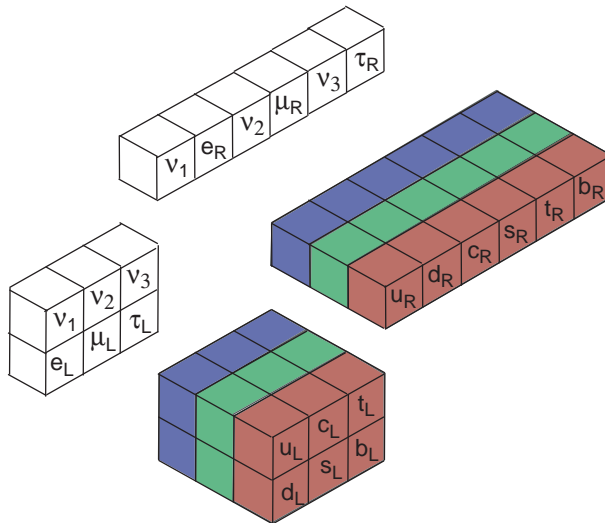
Figure 1. Left-handed doublets and right-handed singlets of quarks and leptons that inspire the structure of the electroweak theory.

energy scale around 1 TeV, or 10^{12} eV, and other considerations also single out the 1-TeV scale as fertile terrain for new physics. Not by coincidence, the Large Hadron Collider will empower experiments to carry out a thorough exploration of the 1-TeV scale. We will describe signatures that will be important in the search for the Higgs boson. Then we will argue, independent of any specific mechanism for electroweak symmetry breaking, that (something like) the Higgs boson must exist. After a brief mention of other new phenomena to be expected on the 1-TeV scale, we will close with a short outlook on the decade of discovery ahead.

2   Sources of Mass in the Electroweak Theory

Our picture of matter is grounded in the recognition of a set of pointlike constituents: the quarks and leptons, as depicted in Figure 1, plus a few fundamental forces derived from gauge symmetries.[1] 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. By pointlike, we understand that the quarks and leptons show no evidence of internal structure at the current limit of our resolution, ($r \lesssim 10^{-18}$ m). It is striking that the charged-current weak interaction responsible for radioactive beta decay and other processes acts only on the left-handed fermions. We do not know whether the observed parity violation reflects a fundamental asymmetry in the laws of Nature, or a hidden symmetry that might be restored at higher energies.

Like quantum chromodynamics, the electroweak theory is a gauge theory, in which interactions follow from symmetries. Let us briefly review the strategy of gauge theories by considering the consequences of local gauge invariance in quantum mechanics. A quantum-mechanical state is described by a complex Schrödinger wave function $\psi(x)$. Quantum-mechanical observables involve inner products of the form

$$\langle \mathcal{O} \rangle = \int d^n x \, \psi^* \mathcal{O} \psi, \tag{2}$$

which are unchanged by a global U(1) phase rotation, $\psi(x) \to e^{i\theta}\psi(x)$. In other words, the absolute phase of the wave function cannot be measured and is a matter of convention. *Relative* phases between wave functions, as measured in interference experiments, are also unaffected by such a global rotation.

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[1]For general surveys of the standard model of particle physics, and a glimpse beyond, see (8; 9).

Are we free to choose independent phase conventions in Batavia and in London? In other words, can quantum mechanics be formulated to remain invariant under local, *position-dependent*, phase rotations, $\psi(x) \to e^{i\alpha(x)}\psi(x)$? It is easy to see that this can be achieved, at the price—or reward—of introducing an interaction that we shall construct to be electromagnetism.

Observables such as momentum, as well as the Schödinger equation itself, involve derivatives of the wave function. Under local phase rotations, these transform as

$$\partial_\mu\psi(x) \to e^{i\alpha(x)}[\partial_\mu\psi(x) + i(\partial_\mu\alpha)\psi(x)], \tag{3}$$

which involves more than a mere phase change. The additional gradient-of-phase term, a four-vector, spoils local phase invariance. But we may achieve local phase invariance if we modify the equations of motion and the definitions of observables by introducing the (four-vector) electromagnetic field $A_\mu(x)$. We replace the normal gradient ∂_μ everywhere—in the Schrödinger equation, definition of the momentum operator, etc.—by the *gauge-covariant derivative* $\mathcal{D}_\mu = \partial_\mu + ieA_\mu$, where e is the charge of the particle described by $\psi(x)$. Then if the field $A_\mu(x)$ transforms under local phase rotations as $A_\mu(x) \to A_\mu(x) - (1/e)\partial\alpha(x)$, local phase rotations take

$$\mathcal{D}_\mu\psi(x) \to e^{i\alpha(x)}\mathcal{D}_\mu\psi(x). \tag{4}$$

Consequently, quantities such as $\psi^*\mathcal{D}_\mu\psi$ are invariant under local phase rotations. The transformation law for the four-vector A_μ has the familiar form of a gauge transformation in electrodynamics. Moreover, the covariant derivative, which prescribes the coupling between matter and the electromagnetic field, corresponds to the replacement $p_\mu \to p_\mu - eA_\mu$ for the momentum operator in the presence of an electromagnetic potential. We have obtained electromagnetism as a consequence of local U(1) phase invariance applied to the Schrödinger wave function.

A parallel strategy can be applied in relativistic quantum field theory for any simple or semi-simple (product) gauge group. The correct electroweak gauge symmetry emerged through trial and error and experimental guidance. To incorporate electromagnetism into a theory of the weak interactions, we add a $U(1)_Y$ weak-hypercharge phase symmetry[1] to the $SU(2)_L$ family (weak-isospin) symmetry suggested by the left-handed doublets of Figure 1. To save writing, we shall display the electroweak theory as it applies to a single generation of leptons. In this form, it is incomplete; for quantum corrections to respect the gauge symmetry, a doublet of color-triplet quarks must accompany each doublet of color-singlet leptons. However, the needed generalizations are simple enough to make that we need not write them out.[2] We begin by specifying the fermions: a left-handed weak isospin doublet

$$\mathsf{L} = \begin{pmatrix} \nu_e \\ e \end{pmatrix}_{\mathrm{L}} \tag{5}$$

with weak hypercharge $Y_{\mathrm{L}} = -1$, and a right-handed weak isospin singlet

$$\mathsf{R} \equiv e_{\mathrm{R}} \tag{6}$$

with weak hypercharge $Y_{\mathrm{R}} = -2$.

The $SU(2)_L \otimes U(1)_Y$ electroweak gauge group implies two sets of gauge fields: a weak isovector $\vec{b}_\mu$, with coupling constant g, and a weak isoscalar $\mathcal{A}_\mu$, with independent coupling constant g'. The gauge fields ensure local gauge invariance, provided they obey the transformation laws $\mathcal{A}_\mu \to \mathcal{A}_\mu - (1/g')\partial_\mu\alpha$ under an infinitesimal hypercharge phase rotation, and $\vec{b}_\mu \to \vec{b}_\mu - \vec{\alpha} \times \vec{b}_\mu - (1/g)\partial_\mu\vec{\alpha}$ under an infinitesimal weak-isospin rotation generated by $G = 1 + (i/2)\vec{\alpha} \cdot \vec{\tau}$, where $\vec{\tau}$ are the Pauli isospin matrices. Corresponding to

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[1]We relate the weak hypercharge Y through the Gell-Mann–Nishijima connection, $Q = I_3 + \frac{1}{2}Y$, to electric charge and (weak) isospin.
[2]The electroweak theory is developed in many textbooks; see especially (10; 11; 12). For a look back at the evolution of the electroweak theory, see the Nobel Lectures by some of its principal architects (13; 14; 15; 16; 17).

these gauge fields are the field-strength tensors

$$F_{\mu\nu}^\ell = \partial_\nu b_\mu^\ell - \partial_\mu b_\nu^\ell + g\varepsilon_{jk\ell} b_\mu^j b_\nu^k \;, \tag{7}$$

for the weak-isospin symmetry, and

$$f_{\mu\nu} = \partial_\nu \mathcal{A}_\mu - \partial_\mu \mathcal{A}_\nu \;, \tag{8}$$

for the weak-hypercharge symmetry.

We may summarize the interactions by the Lagrangian

$$\mathcal{L} = \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{leptons}} \;, \tag{9}$$

with

$$\mathcal{L}_{\text{gauge}} = -\tfrac{1}{4} F_{\mu\nu}^\ell F^{\ell\mu\nu} - \tfrac{1}{4} f_{\mu\nu} f^{\mu\nu}, \tag{10}$$

and

$$\mathcal{L}_{\text{leptons}} = \overline{\mathsf{R}}\, i\gamma^\mu \left(\partial_\mu + i\frac{g'}{2} \mathcal{A}_\mu Y \right) \mathsf{R} \tag{11}$$

$$+ \overline{\mathsf{L}}\, i\gamma^\mu \left(\partial_\mu + i\frac{g'}{2} \mathcal{A}_\mu Y + i\frac{g}{2} \vec{\tau} \cdot \vec{b}_\mu \right) \mathsf{L}.$$

The theory in this form has important shortcomings. The Lagrangian of Eq. 10 contains four massless electroweak gauge bosons, namely $\mathcal{A}_\mu$, b_μ^1, b_μ^2, and b_μ^3, whereas Nature has but one: the photon. (Note that a mass term such as $\frac{1}{2}m^2 \mathcal{A}_\mu \mathcal{A}^\mu$ is not invariant under a gauge transformation.) Moreover, the $\mathrm{SU(2)_L \otimes U(1)_Y}$ gauge symmetry forbids a mass term $m\bar{e}e = m(\bar{e}_\mathrm{R} e_\mathrm{L} + \bar{e}_\mathrm{L} e_\mathrm{R})$ for the electron in Eq. 11, because the left-handed and right-handed fields transform differently. 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 the same symmetry in the outcomes of those laws.

The most apt analogy for the hiding of the electroweak gauge symmetry is found in the superconducting phase transition[1] 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 a nonzero photon mass within the superconducting medium. The Higgs mechanism (21; 22; 23; 24) is a relativistic generalization of the Ginzburg-Landau phenomenology of superconductivity. We introduce a complex doublet of scalar fields

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

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

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

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[1]The parallel between electroweak symmetry breaking and the Ginzburg-Landau theory is drawn carefully in §4.4 of (18). For a rich discussion of superconductivity as a consequence of the spontaneous breaking of electromagnetic gauge symmetry, see §21.6 of (19). For an essay on mass generation through spontaneous symmetry breaking, see (20).

where the gauge-covariant derivative is

$$\mathcal{D}_\mu = \partial_\mu + i\frac{g'}{2}\mathcal{A}_\mu Y + i\frac{g}{2}\vec{\tau}\cdot\vec{b}_\mu , \tag{14}$$

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

$$V(\phi^\dagger\phi) = \mu^2(\phi^\dagger\phi) + |\lambda|\,(\phi^\dagger\phi)^2. \tag{15}$$

We are also free to add a (gauge-invariant) Yukawa interaction between the scalar fields and the leptons,

$$\mathcal{L}_{\text{Yukawa}} = -\zeta_e\left[\overline{\mathsf{R}}(\phi^\dagger\mathsf{L}) + (\overline{\mathsf{L}}\phi)\mathsf{R}\right]. \tag{16}$$

We then arrange their self-interactions so that the vacuum state corresponds to a broken-symmetry solution. The electroweak symmetry is spontaneously broken if the parameter $\mu^2 < 0$. The minimum energy, or vacuum state, may then be chosen to correspond to the vacuum expectation value

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

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

Let us verify that the vacuum of Eq. 17 indeed breaks the gauge symmetry. The vacuum state $\langle\phi\rangle_0$ is invariant under a symmetry operation $\exp(i\alpha\mathcal{G})$ corresponding to the generator $\mathcal{G}$ provided that $\exp(i\alpha\mathcal{G})\langle\phi\rangle_0 = \langle\phi\rangle_0$, *i.e.*, if $\mathcal{G}\langle\phi\rangle_0 = 0$. Direct calculation reveals that the original four generators are all broken, but electric charge is not. We have accomplished our goal of breaking $\mathrm{SU(2)}_{\mathrm{L}}\otimes\mathrm{U(1)}_Y \to \mathrm{U(1)}_{\mathrm{em}}$. 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.

With the definition $g' = g\tan\theta_W$, where θ_W is the weak mixing angle, we can express the photon as the linear combination $A = \mathcal{A}\cos\theta_W + b_3\sin\theta_W$. We identify the strength of its coupling to charged particles, $gg'/\sqrt{g^2 + g'^2}$, with the electric charge e. The mediator of the charged-current weak interaction, $W^\pm = (b_1\mp ib_2)/\sqrt{2}$, acquires a mass $M_W = gv/2 = ev/2\sin\theta_W$. The electroweak gauge theory reproduces the low-energy phenomenology of the Fermi theory of weak interactions, provided we set $v = (G_F\sqrt{2})^{-1/2} = 246$ GeV, where $G_F = 1.16637\times 10^{-5}$ GeV^{-2} is the Fermi constant. It follows at once that $M_W \approx 37.3$ GeV$/\sin\theta_W$. The combination of I_3 and Y 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 $M_Z = M_W/\cos\theta_W$.[1]

As a vestige of the spontaneous symmetry breaking, there remains a massive, spin-zero particle, the Higgs boson. Its mass is given symbolically as $M_H^2 = -2\mu^2 > 0$, but *we have no prediction for its value.* Though what we take to be the work of the Higgs boson is all around us, the Higgs particle itself has not yet been observed!

The masses of the elementary fermions are a more mysterious story: Each fermion mass involves a new, so far incalculable, Yukawa coupling. 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, from $\zeta_e \approx 3\times 10^{-6}$ for the electron to $\zeta_t \approx 1$ for the top quark. Their origin is unknown. In that sense, therefore, *all fermion masses involve physics beyond the standard model.*

Experiments and the supporting theoretical calculations over the past decade have elevated the electroweak theory to a law of Nature, tested as a quantum field theory at the level of one per mille.[2] One remarkable achievement of recent experiments is a clear test of the gauge symmetry, or group-theory

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[1] The weak neutral-current interaction was not known before the electroweak theory. Its discovery in 1973 (25) marked an important milestone, as did the observation a decade later of the $W^\pm$ and Z^0 bosons (26).

[2] The current state of the theory is reviewed in (27). An ongoing comparison of theory and experiment is maintained by the LEP Electroweak Working Group (28).
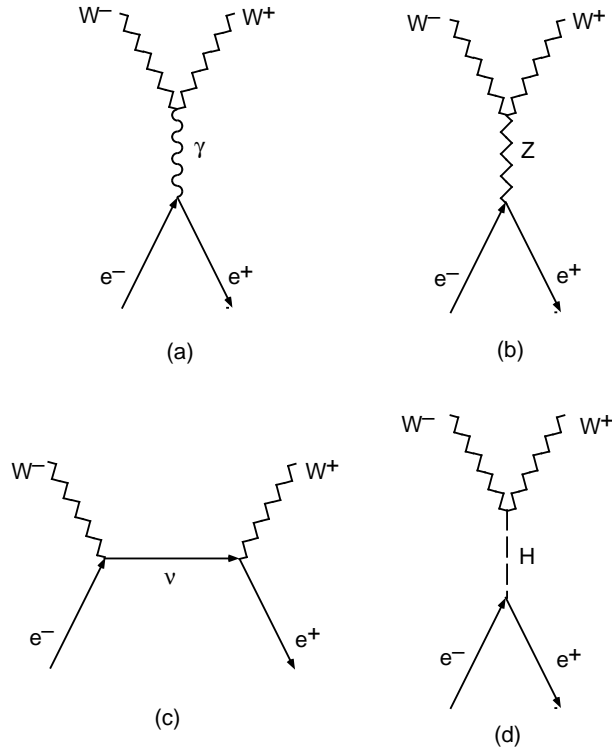
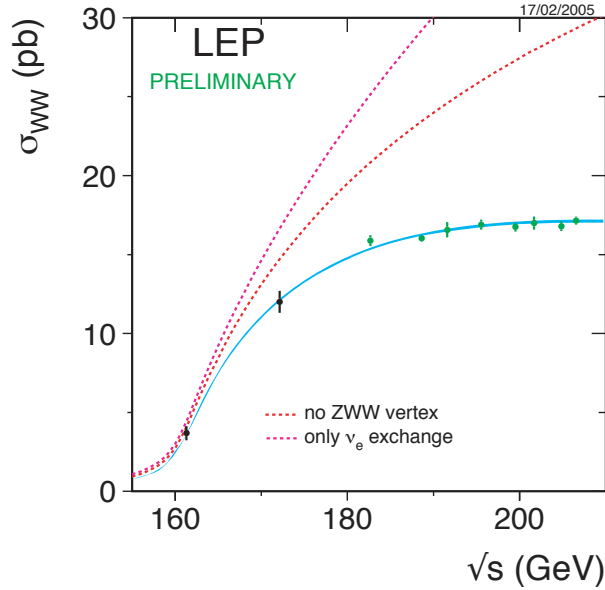Figure 2. Lowest-order contributions to the $e^+e^- \to W^+W^-$ scattering amplitude.

Figure 3. Cross section for the reaction $e^+e^- \to W^+W^-$ measured by the four LEP experiments, together with the full electroweak-theory simulation and the cross sections that would result from ν-exchange alone and from $(\nu + \gamma)$-exchange (28).

structure, of the electroweak theory, in the reaction $e^+e^- \to W^+W^-$. Neglecting the electron mass, this reaction is described by three Feynman diagrams that correspond to s-channel photon and Z^0 exchange, and t-channel neutrino exchange, Figure 2(a-c). Each diagram leads to a $J = 1$ partial-wave amplitude $\propto s$, 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 of Figs. 2(a) and (b) cancel the leading divergence in the $J = 1$ partial-wave amplitude of the neutrino-exchange diagram in Figure 2(c). The LEP measurements (28) plotted in Figure 3 agree well with the benign high-energy behavior predicted

by the electroweak theory. 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—and disagrees with the measurements. The gauge cancellation in the $J = 1$ partial-wave amplitude is thus observed.

3   Why Electroweak Symmetry Breaking Matters to You

Experiments that explore the 1-TeV scale will deepen our understanding of the everyday, the stuff of the world around us, responding in new and revealing ways to the basic questions about atoms, chemistry, and complex objects announced in §1. Perhaps the best way to connect those questions with the electroweak theory and LHC physics is to consider what the world would be like if there were nothing like the Higgs mechanism for electroweak symmetry breaking. First, it's clear that quarks and leptons would remain massless, because mass terms are not permitted if the electroweak symmetry remains manifest.[1] Eliminating the Higgs mechanism does nothing to the strong interaction, so QCD would still confine the (massless) color-triplet quarks into color-singlet hadrons, with very little change in the masses of those stable structures.

An interesting and slightly subtle point is that, even in the absence of a Higgs mechanism, QCD hides the electroweak symmetry (29). In a world with massless up and down quarks, QCD exhibits a global $SU(2)_L \otimes SU(2)_R$ *chiral symmetry* that treats the left-handed and right-handed quarks as separate objects. As we approach low energy from above, that chiral symmetry is spontaneously broken. The resulting communication between the left-handed and right-handed worlds engenders a breaking of the electroweak symmetry: $SU(2)_L \otimes U(1)_Y$ becomes $U(1)_{em}$, and the gauge bosons are the massless photon and massive $W^\pm$ and Z^0. Despite the structural similarity to the standard model, this is not a satisfactory theory of the weak interactions. Here the scale of electroweak symmetry breaking is measured by the pion lifetime—the coupling of the axial current to the vacuum. The amount of mass acquired by the W and Z is too small by a factor of 2500.

Because the weak bosons have acquired masses, the weak-isospin force, which we might have taken to be a confining force in the absence of symmetry breaking, does not confine objects bearing weak isospin into weak-isospin singlets. The familiar spectrum of hadrons persists, but with a crucial difference. The proton, with its electrostatic self-energy, will now outweigh the neutron, because in this world of massless quarks, the down quark now does not outweigh the up quark.

Beta decay—exemplified in this world by $p \to n e^+ \nu_e$—is very rapid, because the gauge bosons are so light. The lightest nucleus is therefore one neutron; *there is no hydrogen atom.* Exploratory analyses of what would happen to big-bang nucleosynthesis in this world suggest that some light elements, such as helium, would be created in the first minutes after the big bang (30; 31; 32; 33). [It would be interesting to see this worked out in complete detail.] Because the electron is massless, the Bohr radius of the atom is infinite, so there is nothing we would recognize as an atom, there is no chemistry as we know it, there are no stable composite structures like the solids and liquids we know.[2]

How very different the world would be, were it not for the mechanism of electroweak symmetry breaking! What we are really trying to get at, when we look for the source of electroweak symmetry breaking, is why we don't live in a world so different, why we live in the world we do. This is one of the deepest questions that human beings have ever tried to engage, and it is coming within the reach of particle physics.

What form might the answer take? What clues we have suggest that the agent of electroweak symmetry breaking represents a novel fundamental interaction at an energy of a few hundred GeV. *We do not know what that force is.*

It could be the Higgs mechanism of the standard model (or a supersymmetric elaboration of the standard model (34)), which is built in analogy to the Ginzburg–Landau description of superconductivity. The

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[1]I assume for this discussion that all the trappings of the Higgs mechanism, including Yukawa couplings for the fermions, are absent.
[2]It is nearly inevitable that effects negligible in our world would, in the Higgsless world, produce fermion masses. These are typically many orders of magnitude smaller than the observed masses, small enough that the Bohr radius of a would-be atom would be macroscopic, sustaining the conclusion that matter would lose its integrity.

potential that we arrange, by decree, to hide the electroweak symmetry arises not from gauge forces but from an entirely new kind of interaction.

Maybe the electroweak symmetry is hidden by a new gauge force. One very appealing possibility—at least until you get into the details—is that the solution to electroweak symmetry breaking will be like the solution to the model for electroweak symmetry breaking, the superconducting phase transition. The superconducting phase transition is first described by the Ginzburg–Landau phenomenology, but then in reality is explained by the Bardeen–Cooper–Schrieffer theory that comes from the gauge theory of Quantum Electrodynamics. Maybe, then, we will discover a mechanism for electroweak symmetry breaking almost as economical as the QCD mechanism we discussed above. One much investigated line is the possibility that new constituents still to be discovered interact by means of yet unknown forces, and when we learn how to calculate the consequences of that theory we will find our analogue of the BCS theory (35).

It could even be that there is some truly emergent description of the electroweak phase transition, a residual force that arises from the strong dynamics among the weak gauge bosons (36). 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 GeV scale, as we shall see in §5. In that event, it is reasonable to speculate that resonances form among pairs of gauge bosons, multiple production of gauge bosons becomes commonplace, and that resonant behavior could hold the key to understanding what hides the electroweak symmetry.

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 "Little Higgs" theories feature weakly coupled new physics at the TeV scale (37).

Or perhaps electroweak symmetry breaking is an echo of extra spacetime dimensions. Among the possibilities are models without a physical Higgs scalar, in which is electroweak symmetry is hidden by boundary conditions (38).

Theory has offered many alternatives. During the next decade, experiment will tell us which path Nature has taken. An essential first step is to find the Higgs boson and to learn its properties. But where shall we look?

4   Higgs-Boson Properties

Once we assume a value for the Higgs-boson mass, it is a simple matter to compute the rates for Higgs-boson decay into pairs of fermions or weak bosons. For a fermion f with N_c colors, the partial width $\Gamma(H \rightarrow f\bar{f})$ is proportional to $N_c m_f^2 M_H$ in the limit of large Higgs mass. The rates for decays into weak-boson pairs are asymptotically proportional to M_H^3 and $\frac{1}{2}M_H^3$, for $H \rightarrow W^+W^-$ and $H \rightarrow ZZ$, respectively, the relative factor $\frac{1}{2}$ arising from weak isospin. 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 4. 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. Though rare, the $\gamma\gamma$ channel offers an important target for LHC experiments, if the Higgs boson is light.

Below the W^+W^- threshold, the total width of the standard-model Higgs boson is rather small, typically less than 1 GeV. Far above the threshold for decay into gauge-boson pairs, the total width is proportional to M_H^3. At masses approaching 1 TeV, the Higgs boson becomes very broad, with a perturbative width approaching its mass. The Higgs-boson total width is plotted as a function of M_H in Figure 5.

As we have described it, the Higgs boson is an artifact of the mechanism we chose to hide the electroweak symmetry. What assurance do we have that a Higgs boson, or something very like it, will be found? One path to the *theoretical* discovery of the Higgs boson involves its role in the cancellation of high-energy divergences. We saw at the end of §2 that the most severe divergences of the individual ν-, γ-, and Z-exchange diagrams in the reaction $e^+e^- \rightarrow W^+W^-$ are tamed by a cooperation among the three
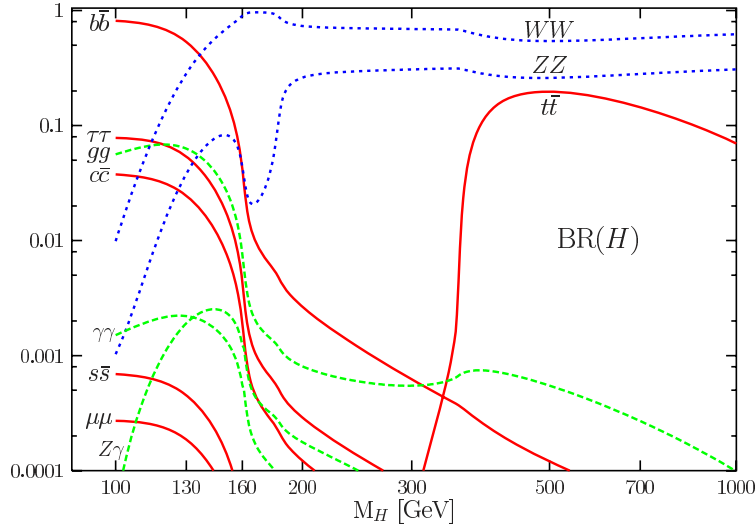
Figure 4. Branching fractions for the prominent decay modes of the standard-model Higgs boson, from (39).
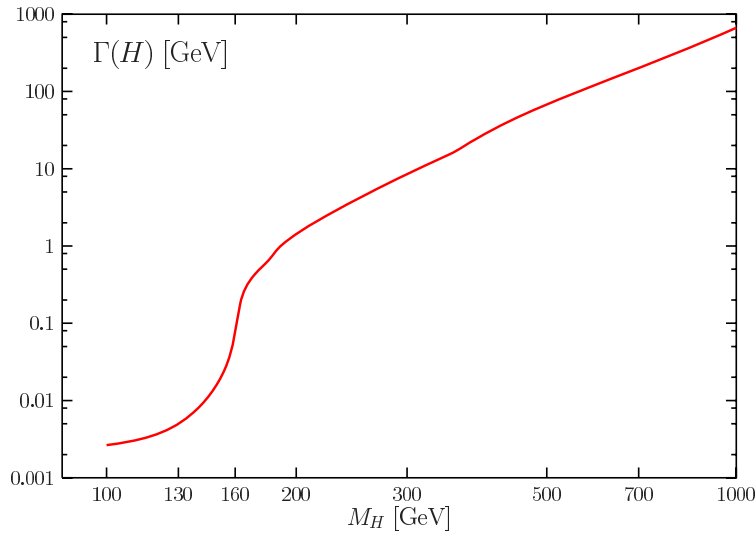

Figure 5. Higgs-boson total width as a function of mass, from (39).

diagrams of Figure 2(a-c) that follows from gauge symmetry. However, this is not the end of the high-energy story: the $J = 0$ partial-wave amplitude, which exists in this case because the electrons are massive and may therefore be found in the "wrong" helicity state, grows as the c.m. energy for the production of longitudinally polarized gauge bosons. This unacceptable high-energy behavior is precisely cancelled by the Higgs boson graph of Figure 2(d). If the Higgs boson did not exist, something else would have to play this role. From the point of view of S-matrix analysis, the Higgs-electron-electron coupling must be proportional to the electron mass, because the strength of "wrong-helicity" configurations is measured by the fermion mass.

Let us underline this result. If the gauge symmetry were unbroken, there would be no Higgs boson, no longitudinal gauge bosons, and no extreme divergence difficulties. But there would be no viable low-energy phenomenology of the weak interactions. The most severe divergences of individual diagrams are eliminated by the gauge structure of the couplings among gauge bosons and leptons. A lesser, but still potentially fatal, divergence arises because the electron has acquired mass—because of the Higgs mechanism. Spontaneous symmetry breaking provides its own cure by supplying a Higgs boson to remove the last divergence. A similar interplay and compensation must exist in any satisfactory theory. There will be (almost surely) a spin-zero object that has effectively more or less the interactions of the standard-model Higgs boson,

whether it be an elementary particle that we build into to the theory or something that emerges from the theory. Such an object is required to make the electroweak theory behave well at high energies, once electroweak symmetry is hidden.

It is by no means guaranteed that the same agent hides electroweak symmetry and generates fermion mass. We saw in §3 that chiral symmetry breaking in QCD could hide the electroweak symmetry without generating fermion masses. In extended technicolor models (40; 41), for example, separate gauge interactions hide the electroweak symmetry and communicate the broken symmetry to the quarks and leptons. 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 4 (42). Accordingly, it will be of great interest to map the decay pattern of the Higgs boson, once it is found, in order to characterize the mechanism of electroweak symmetry breaking.

5   The Importance of the 1-TeV Scale

The electroweak theory does not give a precise prediction for the mass of the Higgs boson, but a unitarity argument (43) leads to a conditional upper bound on the Higgs boson mass that sets a key target for experiment.

It is straightforward to compute the amplitudes $\mathcal{M}$ for gauge boson scattering 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 the particle poles, or at exponentially large energies), for any value of the Higgs boson mass M_H. Four channels are interesting:

$$W_L^+ W_L^- \quad \frac{Z_L^0 Z_L^0}{\sqrt{2}} \quad \frac{HH}{\sqrt{2}} \quad HZ_L^0 \,, \tag{18}$$

where the subscript L denotes the longitudinal polarization states, and 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,[1]

$$\lim_{s \gg M_H^2} (a_0) \to \frac{-G_F M_H^2}{4\pi\sqrt{2}} \cdot \begin{bmatrix} 1 & 1/\sqrt{8} & 1/\sqrt{8} & 0 \\ 1/\sqrt{8} & 3/4 & 1/4 & 0 \\ 1/\sqrt{8} & 1/4 & 3/4 & 0 \\ 0 & 0 & 0 & 1/2 \end{bmatrix} \,. \tag{19}$$

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

$$M_H \leq \left(\frac{8\pi\sqrt{2}}{3G_F}\right)^{1/2} = 1 \text{ TeV} \tag{20}$$

as a condition for perturbative unitarity.

If the bound is respected, weak interactions remain weak at all energies, and perturbation theory is everywhere reliable. If the bound is violated, perturbation theory breaks down, and weak interactions among $W^\pm$, Z, and H become strong on the 1-TeV scale. This means that the features of strong interactions at GeV energies will come to characterize electroweak gauge boson interactions at TeV energies. We interpret this to mean that new phenomena are to be found in the electroweak interactions at energies not much larger than 1 TeV.

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[1]It is convenient to calculate these amplitudes by means of the Goldstone-boson equivalence theorem (44), which reduces the dynamics of longitudinally polarized gauge bosons to a scalar field theory with interaction Lagrangian given by $\mathcal{L}_{\text{int}} = -\lambda v h(2w^+ w^- + z^2 + h^2) - (\lambda/4)(2w^+ w^- + z^2 + h^2)^2$, with $1/v^2 = G_F\sqrt{2}$ and $\lambda = G_F M_H^2/\sqrt{2}$.

If the $SU(2)_L \otimes U(1)_Y$ electroweak theory points to a Higgs boson mass below 1 TeV, it does not explain how the scale of electroweak symmetry breaking is maintained in the presence of quantum corrections. Beyond the classical approximation, scalar mass parameters receive quantum corrections from loops that contain particles of spins $J = 0, \frac{1}{2}$, and 1:
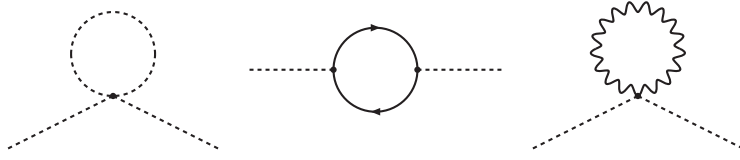
The loop integrals are potentially divergent. Symbolically, we may summarize their implications as

$$M_H^2(p^2) = M_H^2(\Lambda^2) + Cg^2 \int_{p^2}^{\Lambda^2} dk^2 + \cdots , \tag{21}$$

where Λ defines a reference scale at which the value of M_H^2 is known, g is the coupling constant of the theory, and the coefficient C is calculable in any particular theory. Here we describe the variation of an observable with the momentum scale. The loop integrals appear to be quadratically divergent, $\propto \Lambda^2$. In order for the mass shifts induced by quantum corrections to remain under control (*i.e.*, not to greatly exceed the value measured on the laboratory scale), either Λ must be small, so the range of integration is not enormous, or new physics must intervene to damp the integrand.

If the fundamental interactions are described by an $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ gauge symmetry, *i.e.*, by quantum chromodynamics and the electroweak theory, then a natural reference scale is the Planck mass, $\Lambda \sim M_{Planck} = (\hbar c/G_{Newton})^{1/2} \approx 1.2 \times 10^{19}$ GeV. In a unified theory of the strong, weak, and electromagnetic interactions, a natural scale is the unification scale, $\Lambda \sim M_U \approx 10^{15}$ - 10^{16} GeV. Both estimates are very large compared to the electroweak scale. The challenge of preserving widely separated 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 must intervene at an energy of approximately 1 TeV to bring the integral in Eq. 21 under control.

Note the implications: The unitarity argument showed that new physics must be present on the 1-TeV scale, either in the form of a Higgs boson, or other new phenomena. But 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. 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 Eq. 21. We do not have a proof that Nature is not fine tuned, but I think it highly likely that *both a Higgs boson and other new phenomena* are to be found near the 1-TeV scale.

Let us review some of the ways in which new phenomena 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* balances the contributions of fermion and boson loops.[1] In the limit of unbroken supersymmetry, in which the masses of bosons are degenerate with those of their fermion counterparts, the cancellation is exact. If the supersymmetry is broken (as it must be in our world), 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 (46). Theories of dynamical symmetry breaking, such as Technicolor, offer a second solution to the problem of the enormous range of integration in (21). In technicolor models, the Higgs boson is composite, and its internal structure

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[1] "Little Higgs" models and "twin Higgs" models (45) employ different conspiracies of contributions to defer the hierarchy problem to about 10 TeV.

comes into play on the scale of its binding, $\Lambda_{\text{TC}} \simeq \mathcal{O}(1 \text{ TeV})$. The integrand is damped, the effective range of integration is cut off, and mass shifts are under control.

We have one more independent indication that new phenomena should be present on the 1-TeV scale. An appealing interpretation of the evidence that dark matter makes up roughly one-quarter of the energy density of the Universe (47) is that dark matter consists of thermal relics of the big bang, stable—or exceedingly long-lived—neutral particles. If the particle has couplings of weak-interaction strength, then generically the observed dark-matter density results if the mass of the dark-matter particle lies between approximately 100 GeV and 1 TeV (48). Typically, scenarios to extend the electroweak theory and resolve the hierarchy problem—whether based on extra dimensions, new strong dynamics, or supersymmetry—entail dark-matter candidates on the 1-TeV scale. One aspect of the great optimism with which particle physicists contemplate the explorations under way at Fermilab's Tevatron and soon to be greatly extended at CERN's Large Hadron Collider is a strong suspicion that many of the outstanding problems of particle physics and cosmology may be linked—and linked to the 1-TeV scale. Dark matter is a perfect example.

6   Outlook

Over the next decade, experiments will carry out definitive explorations of the Fermi scale, at energies around 1 TeV for collisions among quarks and leptons. This is physics on the *nanonanoscale,* probing distances smaller than 10^{-18} m. In this regime, we confidently expect to find the key to the mechanism that drives electroweak symmetry breaking, with profound implications for our conception of the everyday world. A pivotal step will be the search for the Higgs boson and the elaboration of its properties. What is more, the hierarchy problem leads us to suspect that other new phenomena are to be found on the 1-TeV scale, phenomena that will give new insight into why the electroweak scale is so much smaller than the Planck scale. We also have reason to believe—from arguments about relic densities and also from specific models—that a weakly interacting class of dark-matter candidates could populate the same energy range.

We do not know what the new wave of exploration will find, but the discoveries and new puzzles are certain to change the face of particle physics and reverberate through neighboring disciplines. Resolving the conundrums of the 1-TeV scale should aid us in reformulating some of today's fuzzy questions and give us a clearer view of the physics at still shorter distances, where we may uncover new challenges to our understanding. We could well find new clues to the unification of forces or indications for a rational pattern of constituent masses, viewed at a high energy scale. I hope that we will be able to sharpen the problem of identity—what makes an electron an electron, a top quark a top quark, a neutrino a neutrino—so we can formulate a strategy to resolve it.

Experiments at the Fermilab Tevatron, a 2-TeV proton-antiproton collider, have begun to approach the 1-TeV scale. The CDF (49) and DØ (50) experiments, which discovered the 171-GeV top quark in 1995, are profiting from world-record machine performance: initial luminosities exceeding 2.5×10^{32} cm^2 s^{-1} and an integrated luminosity to date of more than 2 fm^{-1} (51). They are expected to run through 2009, continuing their pursuit of the Higgs boson, supersymmetry, and other new phenomena.

The Large Hadron Collider (52), a 14-TeV proton-proton collider at CERN, will transport us into the heart of the Fermi scale. The LHC should demonstrate 900-GeV collisions by the end of 2007 and produce the first 14-TeV events during 2008.[1] Its collision rate will grow to 100 times the Tevatron's luminosity. Like the machine itself, with its 27-km circumference, the multipurpose detectors ATLAS (54) and CMS (55) are both titans and engineering marvels. They will be the vessels for a remarkable era of exploration and discovery.

Acknowledgments

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[1]For a prospectus on early running at the LHC, see (53).

References

[1] P. A. M. Dirac, Proc. Royal Soc. (London) **A123** 714–733 (1929).

[2] F. Wilczek, Phys. Today **52N11** 11–13 (November 1999).

[3] W. M. Yao *et al.*, J. Phys. **G33** 1–1232 (2006).

[4] S. Aoki *et al.*, Phys. Rev. **D67** 034503 (2003) [`hep-lat/0206009`].

[5] Y. Namekawa *et al.*, Phys. Rev. **D70** 074503 (2004) [`hep-lat/0404014`].

[6] C. Quigg, Phys. Today **50N5** 20–26 (May 1997) [`hep-ph/9704332`].

[7] R. N. Cahn, Rev. Mod. Phys. **68** 951–960 (1996).

[8] M. K. Gaillard, P. D. Grannis and F. J. Sciulli, Rev. Mod. Phys. **71** S96–S111 (1999) [`hep-ph/9812285`].

[9] C. Quigg, Particles and the Standard Model, in G. Fraser (Editor) *The New Physics: for the 21st century* (Cambridge University Press, Cambridge & New York, 2006), chap. 4, pp. 86–118.

[10] I. J. R. Aitchison and A. J. G. Hey, *Gauge Theories in Particle Physics*, third edition , Vol. 2 (Taylor & Francis, London, 2003).

[11] C. Quigg, *Gauge Theories of the Strong, Weak, and Electromagnetic Interactions* (Westview Press, Boulder, Colorado, 1997).

[12] G. D. Coughlan, J. E. Dodd and B. M. Gripaios, *The Ideas of Particle Physics* (Cambridge University Press, Cambridge & New York, 2006).

[13] S. Weinberg, Rev. Mod. Phys. **52** 515–523 (1980).

[14] A. Salam, Rev. Mod. Phys. **52** 525–538 (1980).

[15] S. L. Glashow, Rev. Mod. Phys. **52** 539–543 (1980).

[16] G. 't Hooft, Rev. Mod. Phys. **72** 333–339 (2000).

[17] M. J. G. Veltman, Rev. Mod. Phys. **72** 341–349 (2000).

[18] R. E. Marshak, *Conceptual Foundations of Modern Particle Physics* (World Scientific, Singapore, 1993).

[19] S. Weinberg, *The Quantum Theory of Fields*, Vol. 2 (Cambridge University Press, Cambridge, 1996).

[20] F. Wilczek, Phys. Today **53N1** 13–14 (January 2000).

[21] P. W. Higgs, Phys. Lett. **12** 132–133 (1964).

[22] F. Englert and R. Brout, Phys. Rev. Lett. **13** 321–322 (1964).

[23] P. W. Higgs, Phys. Rev. Lett. **13** 508–509 (1964).

[24] P. W. Higgs, Phys. Rev. **145** 1156–1163 (1966).

[25] D. Haidt, Eur. Phys. J. **C34** 25–31 (2004).

[26] C. Rubbia, Rev. Mod. Phys. **57** 699–722 (1985).

[27] S. de Jong, PoS **HEP2005** 397 (2006) [`hep-ex/0512043`].

[28] The LEP Electroweak Working Group [`lepewwg.web.cern.ch`].

[29] M. Weinstein, Phys. Rev. **D8** 2511 (1973).

[30] V. Agrawal, S. M. Barr, J. F. Donoghue, *et al.*, Phys. Rev. Lett. **80** 1822–1825 (1998).

[31] V. Agrawal, S. M. Barr, J. F. Donoghue, *et al.*, Phys. Rev. D **57** 5480–5492 (1998).

[32] C. J. Hogan, Rev. Mod. Phys. **72** 1149–1161 (2000).

[33] J. J. Yoo and R. J. Scherrer, Phys. Rev. D **67** 043517 (2003).

[34] S. P. Martin, A supersymmetry primer. fourth edition (2006) [`hep-ph/9709356v4`].

[35] C. T. Hill and E. H. Simmons, Phys. Rept. **381** 235–402 (2003) [`hep-ph/0203079`] [Erratum-ibid. **390** 553-554 (2004)].

[36] M. S. Chanowitz, Czech. J. Phys. **55** B45–B58 (2005) [`hep-ph/0412203`].

[37] M. Schmaltz and D. Tucker-Smith, Ann. Rev. Nucl. Part. Sci. **55** 229–270 (2005) [`hep-ph/0502182`].

[38] C. Csaki, J. Hubisz and P. Meade, Electroweak symmetry breaking from extra dimensions, in J. Terning, C. E. M. Wagner and D. Zeppenfeld (Eds) *Physics in D ≥ 4: TASI 2004* (World Scientific, Singapore, 2006), pp. 703–776 [`hep-ph/0510275`].

[39] A. Djouadi, The anatomy of electro-weak symmetry breaking. I: The Higgs boson in the standard model. (2005) [`hep-ph/0503172`].

[40] S. Dimopoulos and L. Susskind, Nucl. Phys. **B155** 237–252 (1979).

[41] E. Eichten and K. D. Lane, Phys. Lett. **B90** 125–130 (1980).

[42] A. Djouadi, The anatomy of electro-weak symmetry breaking. II: The Higgs bosons in the minimal supersymmetric model. (2005) [`hep-ph/0503173`].

[43] B. W. Lee, C. Quigg and H. B. Thacker, Phys. Rev. **D16** 1519 (1977).

[44] J. M. Cornwall, D. N. Levin and G. Tiktopoulos, Phys. Rev. **D10** 1145 (1974) [Erratum-ibid. **D11** 972 (1975)].

[45] Z. Chacko, H. S. Goh and R. Harnik, Phys. Rev. Lett. **96** 231802 (2006) [`hep-ph/0506256`].

[46] U. Amaldi, W. de Boer and H. Fürstenau, Phys. Lett. **B260** 447–455 (1991).

[47] D. N. Spergel *et al.*, Wilkinson Microwave Anisotropy Probe (WMAP) three year results: Implications for cosmology. (2006) [`astro-ph/0603449`].

[48] G. Bertone, D. Hooper and J. Silk, Phys. Rept. **405** 279–390 (2005) [`hep-ph/0404175`].

[49] The Collider Detector at Fermilab [`www-cdf.fnal.gov`].

[50] The DØ Experiment [`www-d0.fnal.gov`].

[51] Tevatron Luminosity Charts [`www.fnal.gov/pub/now/tevlum.html`].

[52] The Large Hadron Collider Project [`lhc.web.cern.ch/lhc`].

[53] F. Gianotti and M. L. Mangano, LHC physics: The first one-two year(s), in G. Carlino and P. Paolucci (Eds) *Proceedings of the 2nd Italian Workshop on the Physics of Atlas and CMS*, Frascati physics series 38 (INFN, Frascati, 2005), pp. 3–26 [`hep-ph/0504221`].

[54] The ATLAS Experiment [`atlasexperiment.org`].

[55] Compact Muon Solenoid [`cms.cern.ch`].