Electroweak Symmetry Breaking in Historical Perspective

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Abstract

The discovery of the Higgs boson is a major milestone in our progress toward understanding the natural world. A particular aim of this review is to show how diverse ideas came together in the conception of electroweak symmetry breaking that led up to the discovery. I also survey what we know now that we did not know before, what properties of the Higgs boson remain to be established, and what new questions we may now hope to address.

Contents

1. INTRODUCTION

A lively continuing conversation between experiment and theory has brought us to a radically simple conception of the material world. Fundamental particles called quarks and leptons are the stuff of direct experience, and two new laws of nature govern their interactions. Pursuing clues from experiment, theorists have constructed the electroweak theory (1–3) and quantum chromodynamics (QCD) (4–7), refined them within the framework of local gauge symmetries, and elaborated their consequences. In the electroweak theory, electromagnetism and the weak interactions—so different in range and apparent strength—are ascribed to a common gauge symmetry. We say that the electroweak gauge symmetry is broken, by dynamics or circumstances, to the gauge symmetry of electromagnetism.

Quark: an elementary spin-1/2 particle that experiences the strong interaction; the current roster is composed of six species, grouped in three weak-isospin doublets

Lepton: an elementary spin-1/2 particle that does not experience the strong interaction; the current roster is composed of three charged particles and three neutrinos

Higgs boson: an elementary scalar particle that is the avatar of electroweak symmetry breaking in the standard electroweak theory, an excitation of the auxiliary scalar fields introduced to contrive a vacuum that does not respect the full $SU(2)_L \otimes U(1)_Y$ symmetry on which the electroweak theory is built

The electroweak theory and QCD join to form the Standard Model of particle physics. Augmented to incorporate neutrino masses and lepton mixing, the Standard Model describes a vast array of experimental information. Experiments have validated the gauge theories of the strong, weak, and electromagnetic interactions to an extraordinary degree as relativistic quantum field theories. Recent textbook treatments of QCD and the electroweak theory may be found in, for example, References 8–11.

Until recently, the triumph of this new picture has been incomplete, notably because we had not identified the agent that differentiates electromagnetism from the weak interaction. The 2012 discovery of the Higgs boson by the ATLAS (12) and CMS (13) Collaborations working at CERN's Large Hadron Collider (LHC) capped a four-decade-long quest for that agent. (Further details of the discoveries are reported in References 14–17.) The observations indicate that the electroweak symmetry is spontaneously broken, or hidden: The vacuum state does not exhibit the full symmetry on which the theory is founded. Crucial insights into spontaneously broken gauge theories were developed a half-century ago by Englert & Brout (18), Higgs (19, 20), and Guralnik et al. (21). All the experimental information we have (22, 23)[1] tells us that the unstable 125-GeV particle discovered by the LHC experiments behaves as an elementary scalar, consistent with the properties anticipated for the Standard Model Higgs boson.

The first goal of this review is to describe how a broad range of concepts, drawn mainly from weak-interaction phenomenology, gauge field theories, and condensed matter physics, came together in the electroweak theory. The presentation complements the construction of the

[1] Public results from the CMS Collaboration pertaining to the Higgs boson may be found at **http://j.mp/1BriIBt**; those from the ATLAS Collaboration may be found at **http://j.mp/1A1kzA5**.

electroweak theory given in my prediscovery article, "Unanswered Questions in the Electroweak Theory" (24). Presentations that are similar in spirit may be found in References 25 and 26. Next, I briefly summarize what we now know about the Higgs boson, what the discovery has taught us, and why the discovery is important to our conception of nature. Finally, I address what remains to be learned about the 125-GeV Higgs boson and what new questions are raised by its existence. For example, we need to discover what accounts for the masses of the electron and the other leptons and quarks, without which there would be no atoms, no chemistry, and no liquids or solids—no stable structures. In the standard electroweak theory, both tasks are the work of the Higgs boson. Moreover, we have reason to believe that the electroweak theory is imperfect, and that new symmetries or new dynamical principles are required to make it fully robust. Throughout this review, I emphasize concepts over technical details.

ATLAS: one of two general-purpose experiments for the Large Hadron Collider (LHC), located adjacent to CERN's main campus

CMS: the Compact Muon Solenoid, one of two general-purpose experiments for the LHC; located in Cessy, France

CERN: the European Laboratory for Particle Physics, which 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 21 member states

LHC: the Large Hadron Collider at CERN is a two-bore proton synchrotron with a circumference of 27 km; it is designed to provide proton–proton collisions up to a center-of-mass energy of 14 TeV and a luminosity exceeding 10^{34} cm $^{-2}$ s $^{-1}$, as well as Pb–Pb and proton–Pb collisions

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

2. EXPERIMENTAL ROOTS OF THE ELECTROWEAK THEORY

In order to establish what a successful theory would need to explain, I devote this section to a compressed evocation of how the phenomenology of the (charged-current) weak interactions developed. A superb reference for the experimental observations that led to the creation of the Standard Model is the book by Cahn & Goldhaber (27), which discusses and reproduces many classic papers.

Becquerel's (28) discovery of radioactivity in 1896 is one of the wellsprings of modern physics. In a short time, physicists learned to distinguish several sorts of radioactivity, classified by Rutherford (29) according to the character of the energetic projectile emitted in the spontaneous disintegration. Natural and artificial radioactivity includes nuclear β decay, observed as

$$^{\mathrm{A}}\mathrm{Z} \rightarrow {}^{\mathrm{A}}(\mathrm{Z}+1) + \beta^{-}, \qquad 1.$$

where β^{-} is Rutherford's name for what was soon identified as the electron and $^{\mathrm{A}}\mathrm{Z}$ stands for the nucleus with charge Z and mass number A (in modern language, Z protons and A-Z neutrons). Examples are tritium β decay, $^{3}\mathrm{H}_1 \rightarrow {}^{3}\mathrm{He}_2 + \beta^{-}$; neutron β decay, $n \rightarrow p + \beta^{-}$; and β decay of lead-214, $^{214}\mathrm{Pb}_{82} \rightarrow {}^{214}\mathrm{Bi}_{83} + \beta^{-}$.

For two-body decays, as indicated by the detected products, the Principle of Conservation of Energy and Momentum states that the β particle should have a definite energy. What was observed, as experiments matured, was very different: In 1914, Chadwick (30), later to discover the neutron, showed conclusively that in the decay of radium B and C (^{214}Pb and ^{214}Bi), the β energy follows a continuous spectrum.

The β-decay energy crisis tormented physicists for years. On December 4, 1930, Pauli (31) addressed an open letter to a meeting on radioactivity in Tübingen, Germany. In his letter, Pauli advanced the outlandish idea of a new, very penetrating, neutral particle of vanishingly small mass. Because Pauli's new particle interacted very feebly with matter, it would escape undetected from any known apparatus, taking with it some energy, which would seemingly be lost. The balance of energy and momentum would be restored by the particle we now know as the electron's antineutrino. Accordingly, the proper scheme for β decay is

$$^{\mathrm{A}}\mathrm{Z} \rightarrow {}^{\mathrm{A}}(\mathrm{Z}+1) + \beta^{-} + \bar{\nu}. \qquad 2.$$

What Pauli called his "desperate remedy" was, in its way, very conservative, for it preserved the principle of energy and momentum conservation and with it the notion that the laws of physics are invariant under translations in space and time.

After Chadwick's (32) discovery of the neutron in highly penetrating radiation emitted by beryllium irradiated by α particles, Fermi (33) named Pauli's hypothetical particle the neutrino, so as to distinguish it from Chadwick's strongly interacting neutron, and constructed his four-fermion theory (what we today call a low-energy effective field theory) of β decay, which was the first step toward the modern theory of the charged-current weak interaction. In retrospect, nuclear β decay was the first hint of flavor, the existence of particle families containing distinct species. That hint was made manifest by the discovery of the neutron, nearly degenerate in mass with the proton, which suggested that the neutron and proton might be two states of a nucleon; the mass difference between the two was attributed to electromagnetic effects. The inference that the neutron and proton were partners was strengthened by the observation that nuclear forces are charge independent up to electromagnetic corrections (34). The accumulating evidence inspired Heisenberg (35) and Wigner (36) to make an analogy between the proton and neutron on the one hand and the up and down spin states of an electron on the other. Isospin symmetry, based on the spin-symmetry group SU(2), is the first example of a flavor symmetry.

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 be a valid description of physics at energies only below the masses of the heavy particles, and must be superseded by a more complete (but perhaps still effective) theory on that energy scale

Detecting a particle that interacts as feebly as the neutrino requires a massive target and a copious source of neutrinos. In 1953, Reines & Cowan (37) used the intense flux of antineutrinos from a fission reactor and a heavy target (10.7 feet^3 of liquid scintillator) containing approximately 10^{28} protons to detect the inverse neutron-β-decay reaction $\bar{\nu} + p \to e^{+}+n$. Initial runs at the Hanford Engineering Works in Benton County, Washington, were suggestive but inconclusive. Moving their apparatus to the stronger fission neutrino source at the Savannah River nuclear plant in South Carolina, Cowan, Reines, and their team (38) made the definitive observation of inverse β decay in 1956.

Through the 1950s, a series of experimental puzzles led to the suggestion that the weak interactions did not respect reflection symmetry, or parity (39). In 1956, Wu and collaborators (40) detected a correlation between the spin vector $\vec{J}$ of a polarized ^{60}Co nucleus and the direction $\hat{p}_e$ of the outgoing β particle. Now, parity inversion leaves spin, an axial vector, unchanged ($P{:}\vec{J} \to \vec{J}$), while reversing the electron direction ($P{:}\hat{p}_e \to -\hat{p}_e$), so the correlation $\vec{J}\cdot\hat{p}_e$ should be an "unobservable" null quantity if parity is a good symmetry. The observed correlation is parity violating. A detailed analysis of the ^{60}Co result, and others that came out in quick succession, established that the charged-current weak interactions are left-handed. By the same argument, the parity operation links a left-handed neutrino with a right-handed neutrino. Therefore, a theory that contains only ν_{L} would be manifestly parity violating.

Could the neutrino indeed be left-handed? Goldhaber and collaborators (41) inferred the electron neutrino's helicity from the longitudinal polarization of the recoil nucleus in the electron-capture reaction

$$\begin{aligned} e^- + {}^{152}\mathrm{Eu}^m(J=0) \to {}^{152}&\mathrm{Sm}^*(J=1) + \nu_e \\ &\hookrightarrow \gamma + {}^{152}\mathrm{Sm}. \end{aligned} \qquad 3.$$

A compendious knowledge of the properties of nuclear levels, together with meticulous technique, enabled this classic experiment.

Following the observation of maximal parity violation in the late 1950s, one could write a serviceable effective Lagrangian for the weak interactions of electrons and neutrinos as the product of charged leptonic currents:

$$\mathcal{L}_{\mathrm{V-A}} = \frac{-G_{\mathrm{F}}}{\sqrt{2}}\bar{\nu}\gamma_\mu(1-\gamma_5)e\ \bar{e}\gamma^\mu(1-\gamma_5)\nu + \text{h.c.}, \qquad 4.$$

where Fermi's coupling constant, G_{F}, is $1.1663787(6)\times 10^{-5}\,\mathrm{GeV}^{-2}$. This Lagrangian has a V − A (vector minus axial vector) Lorentz structure (42–45), whereas Fermi's effective Lagrangian for β decay was a (parity-conserving) vector interaction. A straightforward lepton current–times–nucleon current generalization of Equation 4 that accounts for the fact that nucleons are not simple Dirac particles leads to an effective Lagrangian for β decay and associated processes. Many applications and experimental tests are described in detail elsewhere (46–48).

The direct phenomenological consequences of parity violation in the weak interactions, which shattered the received wisdom of the era, were themselves dramatic. For example, they led to a factor-of-three difference between the total cross sections for ν_e and $\bar{\nu}_e$ scattering. Parity violation is also a harbinger of a particular challenge to be met by a true theory of the weak interactions. In quantum electrodynamics (QED), it is perfectly respectable (and correct!) to write a Lagrangian that includes a term for electron mass,

$$\mathcal{L} = \bar{e}(i\gamma^\mu \mathcal{D}_\mu - m)e = \bar{e}(i\gamma^\mu \partial_\mu - m)e - qA_\mu \bar{e}\gamma^\mu e, \qquad 5.$$

where A_μ is the four-vector potential of electromagnetism. The left-handed and right-handed components of the electron have the same charge, so they appear symmetrically. If fermions are chiral, which is to say that the left-handed and right-handed components behave differently, then a mass term conflicts with symmetries. Section 3.2 describes this issue more precisely.

A second charged lepton, the muon, was discovered and identified as lacking strong interactions in the decade beginning in 1937 (49–51). Similar to the electron, the muon is a spin-1/2 Dirac particle, structureless at our present limits of resolution. It is unstable, with a mean lifetime of approximately 2.2 μs and a mass 207 times that of the electron. It might be tempting, therefore, to consider the muon an excited electron, but the transitions $\mu \to e\gamma$, $\mu \to ee^+e^-$, and $\mu \to e\gamma\gamma$ have never been observed. The limits on these decays are so stringent [e.g., the branching fraction for $\mu \to e\gamma$ is $< 5 \times 10^{-13}$ at the 90% confidence level (22)] that we consider the muon a distinct lepton species.

If the muon is distinct from the electron, what is the nature of the neutrino produced in association with the muon in pion decay, $\pi^+ \to \mu^+\nu$? In 1962, Lederman, Schwartz, Steinberger, and their collaborators carried out a two-neutrino experiment using neutrinos created in the decay of high-energy pions from the new Alternating Gradient Synchrotron at Brookhaven National Laboratory, New York (52). They observed numerous examples of the reaction $\nu N \to \mu + X$ but found no evidence of the production of electrons. Their study established that the muon produced in pion decay is a distinct particle, ν_μ, that is different from both ν_e and $\bar{\nu}_e$. This observation suggests that the leptonic charged-current weak interactions exhibit a two-doublet family structure,

$$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_{\mathrm{L}} \qquad \begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_{\mathrm{L}}. \qquad 6.$$

We are led to generalize the effective Lagrangian of Equation 4 to include the terms

$$\mathcal{L}^{(e\mu)}_{\mathrm{V-A}} = \frac{-G_{\mathrm{F}}}{\sqrt{2}} \bar{\nu}_\mu \gamma_\mu (1-\gamma_5)\mu \; \bar{e}\gamma^\mu(1-\gamma_5)\nu_e + \mathrm{h.c.} \qquad 7.$$

in the familiar current–current form.

Because the weak interaction acts at a point, the effective Lagrangians hold only over a finite range of energies, and they cannot reliably be computed beyond leading order. A classic application (53) of partial-wave unitarity (probability conservation) to inverse muon decay, $\nu_\mu e \to \mu \nu_e$, leads to the conclusion that the four-fermion effective Lagrangian of Equation 7 can make sense only for center-of-mass (c.m.) energies $\sqrt{s} \le 617$ GeV. That comfortably encompasses most laboratory experiments but as a matter of principle gives a clear lesson: New physics must intervene below a c.m. energy of approximately 600 GeV.

Although Fermi took his inspiration from the theory of electromagnetism, he did not posit a force carrier analogous to the photon. This is a perfectly reasonable first step, given that electromagnetism acts over an infinite range, whereas the influence of the β decay interaction extends only over approximately 10^{-15} cm. One may hope to obtain a more satisfactory theory by taking the next step, supposing that the weak interaction, like QED, is mediated by vector boson exchange (of nonzero range) to soften the high-energy growth of amplitudes. The weak intermediate boson must carry charge ± 1, because the familiar manifestations of the weak interactions (such as β decay) are charge changing; they must be rather massive ($\sim$100 GeV), to reproduce the short range of the weak force; and they must accommodate parity violation. Introducing a weak boson $W^{\pm}$ in this ad hoc manner indeed mitigates the unitarity problem for inverse muon decay but introduces incurable unitarity problems for reactions such as $e^+e^- \rightarrow W^+W^-$ or $\nu\bar{\nu} \rightarrow W^+W^-$, as detailed in section 6.2 of Reference 9.

Superconductivity: a phenomenon that occurs in many materials when they are cooled to low temperatures or subjected to high pressure; entails zero electrical resistance and the expulsion of magnetic fields

It is also worth mentioning the discovery of strange particles in the early 1950s, because it was essential to establishing that the leptonic and hadronic weak interactions have the same strength and stimulated the invention of quarks (54–56). Semileptonic decays of hyperons (57) were an essential testing ground for Cabibbo's (58) formulation of the universality of the charged-current weak interactions, which was the forerunner of today's 3×3 quark-mixing matrix (59).

3. THE DEVELOPMENT OF THE ELECTROWEAK THEORY

This section is a brief historical survey of the ideas that came together in the concept of a gauge theory for the weak and electromagnetic interactions. What follows is neither a complete intellectual history (which would occupy a book) nor an abbreviated course, but rather a tour of key themes, including Yang–Mills theory, the insight from superconductivity that spontaneous breaking of a gauge symmetry endows gauge bosons with mass, and the development of the electroweak theory as we know it. The aims are to stress the interplay among ideas from diverse sources and to show how the electroweak theory responds to the established phenomenology of the weak interactions.

3.1. Symmetries and Interactions

Notions of symmetry lie at the heart of much of science, and a confidence in the importance of symmetry is a guiding principle for scientists in many disciplines. Heisenberg's (60, p. 280) quasi-Biblical pronouncement, »*Am Anfang war die Symmetrie*« ("In the beginning was symmetry"),[2] resonates in much theoretical work from the early twentieth century to the present. An essential insight of our modern conception of nature is that symmetries dictate interactions.

Although Heisenberg's assertion can be challenged as mere opinion, physicists have learned over the past century how to connect symmetries with conservation laws, and symmetries with interactions. In 1918, two mathematical theorems by Noether (61, 62) showed that for every continuous global symmetry of the laws of nature there exists a corresponding conservation law. Thus, translation invariance in space—the statement that the laws are the same everywhere—is connected with conservation of momentum. Invariance under translations in time is correlated with the conservation of energy. Invariance under rotations implies the conservation of angular momentum. Noether's theorems show how conservation laws could arise and, indeed, how they could be exact statements, not merely summaries of empirical evidence.

[2] *»Am Anfang war die Symmetrie«, das ist sicher richtiger als die Demokritsche These »Am Anfang war das Teilchen«. Die Elementarteilchen verkörpern die Symmetrien, sie sind ihre einfachsten Darstellungen, aber sie sind erst eine Folge der Symmetrien.* ("In the beginning was Symmetry," that is surely more correct than the Democritean thesis, "In the beginning was the particle." The elementary particles embody symmetries, they are their simplest representations, but they are above all a consequence of the symmetries.)

The derivation of interactions from symmetries was initiated by Weyl in a series of papers, published from 1918 to 1929, spanning the invention of quantum mechanics (63–66).[3] In the version that became a prototype for modern gauge theories, Weyl showed that by requiring that the laws of nature be invariant under local changes of the phase convention for the quantum-mechanical wave function, $\psi(x) \to \psi'(x) = e^{i\alpha(x)}\psi(x)$, one can derive the laws of electrodynamics. Invariance under global (coordinate-independent) $\mathrm{U(1)_{em}}$ phase rotations implies the conservation of electric charge; invariance under local (coordinate-dependent) $\mathrm{U(1)_{em}}$ phase rotations implies the existence of a massless vector field—the photon—that couples minimally to the conserved current of the theory. A straightforward derivation leads to the Lagrangian

$$\begin{aligned}\mathcal{L}_{\mathrm{QED}} &= \mathcal{L}_{\mathrm{free}} - J^{\mu}A_{\mu} - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} \\ &= \bar{\psi}(i\gamma^{\mu}\mathcal{D}_{\mu} - m)\psi - \frac{1}{4}F_{\mu\nu}F^{\mu\nu},\end{aligned} \qquad 8.$$

where ψ is the electron field, $\partial_\mu + iqA_\mu(x) \equiv \mathcal{D}_\mu$ is the gauge-covariant derivative, $J^\mu = q\bar{\psi}\gamma^\mu\psi$ is the conserved electromagnetic current, and the field-strength tensor is $F^{\mu\nu} = -F^{\nu\mu} = \partial^\nu A^\mu - \partial^\mu A^\nu$. The $F_{\mu\nu}F^{\mu\nu}$ term, which accounts for photon propagation, is called the kinetic term. Under a local phase rotation, the photon field transforms as $A_\mu(x) \to A_\mu(x) - \partial_\mu\alpha(x)$, the familiar form of a gauge transformation in (even classical) electrodynamics. The electron mass term $(-m\bar{\psi}\psi)$ respects the local gauge symmetry. A photon mass term would have the form $L_\gamma = \frac{1}{2}m^2A^\mu A_\mu$, which conflicts with local gauge invariance because $A^\mu A_\mu \to (A^\mu - \partial^\mu\alpha)(A_\mu - \partial_\mu\alpha) \neq A^\mu A_\mu$. Thus has local gauge invariance led to the existence of a massless photon.

The construction of QED as the gauge theory (67) based on $\mathrm{U(1)_{em}}$ phase symmetry provides a template for building other interactions derived from symmetries. In 1954, as isospin emerged as a reliable classification symmetry for nuclear levels and as a tool for understanding nuclear forces, Yang & Mills [68; also see Shaw (69)] asked whether isospin, promoted to a local symmetry, could lead to a theory of nuclear forces. It is a lovely idea: Derive the strong interactions among nucleons by requiring that the theory be invariant under independent choices at every point of the convention defining proton and neutron.

The construction begins with the free-nucleon Lagrangian

$$\mathcal{L}_0 = \bar{\psi}(i\gamma^\mu\partial_\mu - m)\psi, \qquad 9.$$

written in terms of the composite fermion fields $\psi \equiv \binom{p}{n}$. This Lagrangian is invariant under global isospin rotations $\psi \to \exp(i\boldsymbol{\tau}\cdot\boldsymbol{\alpha}/2)\,\psi$, where $\boldsymbol{\tau}$ is a Pauli isospin matrix, and the isospin current $J^\mu = \bar{\psi}\gamma^\mu\frac{\boldsymbol{\tau}}{2}\psi$ is conserved. Now require invariance under a local gauge transformation, $\psi(x) \to \psi'(x) = G(x)\psi(x)$, where $G(x) \equiv \exp[i\boldsymbol{\tau}\cdot\boldsymbol{\alpha}(x)/2]$. The construction is similar to the one made for QED, but is more involved because of the non-Abelian nature of the SU(2) isospin gauge group. In this case, we find an isovector of gauge fields, corresponding to the adjoint representation of SU(2). The gauge fields satisfy the transformation law $\mathbf{b}'_\mu = \mathbf{b}_\mu - \boldsymbol{\alpha}\times\mathbf{b}_\mu - (1/g)\partial_\mu\boldsymbol{\alpha}$ or, in component form, $b'^l_\mu = b^l_\mu - \varepsilon_{jkl}\alpha^j b^k - (1/g)\partial_\mu\alpha^l$—the translation familiar from QED plus an isospin rotation. Here, g is the coupling constant of the theory. The field-strength tensor is $F^l_{\mu\nu} = \partial_\nu b^l_\mu - \partial_\mu b^l_\nu + g\varepsilon_{jkl}b^j_\mu b^k_\nu$. It is convenient to define $F_{\mu\nu} = \frac{1}{2}F^l_{\mu\nu}\tau^l$. Then, we may write the

[3]In Reference 63, the symmetry is called *Maßstab-Invarianz*. The terminology *Eichinvarianz* enters in Reference 64 and in Reference 65, chapter IV, section 35, p. 282. The ultimate formulation is in Reference 66.

Yang–Mills Lagrangian as

$$
\begin{aligned}
\mathcal{L}_{\mathrm{YM}} &= \bar{\psi}(i\gamma^{\mu}\partial_{\mu} - m)\psi - \frac{g}{2}\mathbf{b}_{\mu}\cdot\bar{\psi}\gamma^{\mu}\boldsymbol{\tau}\psi - \frac{1}{2}\mathrm{tr}\left(F_{\mu\nu}F^{\mu\nu}\right) \\
&= \mathcal{L}_0 - \frac{g}{2}\mathbf{b}_{\mu}\cdot\bar{\psi}\gamma^{\mu}\boldsymbol{\tau}\psi - \frac{1}{2}\mathrm{tr}\left(F_{\mu\nu}F^{\mu\nu}\right),
\end{aligned} \qquad 10.
$$

which is a free Dirac Lagrangian plus an interaction term that couples the isovector gauge fields to the conserved isospin current, plus a kinetic term that now describes both the propagation and the self-interactions of the gauge fields. As in the case of electromagnetism, a mass term that is quadratic in the gauge fields is incompatible with local gauge invariance, but nothing forbids a common nonzero mass for the nucleons. The quadratic term in the gauge fields present in the field-strength tensor gives rises to self-interactions among the gauge bosons that are not present in Abelian theories such as QED.

The discovery that interactions may be derived from isospin symmetry, and from a general gauge group (70), provides theorists with an important strategy for deriving potentially well-behaved theories of the fundamental interactions. Nuclear forces are not mediated by massless spin-1 particles, so the Yang–Mills theory does not succeed in the goal that motivated it. Nevertheless, the approach underlies two new laws of nature: QCD and the electroweak theory.

3.2. $\mathrm{SU(2)_L \otimes U(1)}_Y$

The Yang–Mills experience shows that there is no guarantee that a gauge theory built on a particular symmetry will faithfully describe some aspect of matter. A great deal of art and, to be sure, trial and error goes into the selection of the right gauge symmetry. In the late 1950s and early 1960s, several authors advanced proposals for a gauge theory of the weak interactions, or of a unified theory of the weak and electromagnetic interactions, reading clues from experiment as best they could. Even after what would become the standard $\mathrm{SU(2)_L \otimes U(1)}_Y$ electroweak theory had emerged and was elaborated, imaginative theorists put forward alternative ideas, guided either by experimental hints or by aesthetics. We do not (yet) have a way of deducing the correct gauge symmetry from higher principles.

What turned out to be the correct choice was elaborated by Glashow (1) in 1961. Let us review the essential structure to recall why a new idea was needed to arrive at a successful theory, even after the correct symmetry had been chosen. The leptonic elements of the theory suffice to exhibit the motivation and the principal features.

We begin by designating the spectrum of fundamental fermions of the theory. It suffices for the moment to include only the electron and its neutrino, which form a left-handed "weak-isospin" doublet (see Equation 6), $\mathsf{L}_e \equiv \binom{\nu}{e}_{\mathrm{L}}$, where the left-handed states are $\nu_{\mathrm{L}} = \frac{1}{2}(1-\gamma_5)\nu$ and $e_{\mathrm{L}} = \frac{1}{2}(1-\gamma_5)e$. For the reasons reviewed in Section 2, it is convenient to assume that the right-handed state $\nu_{\mathrm{R}} = \frac{1}{2}(1+\gamma_5)\nu$ does not exist. Thus, we designate only one right-handed lepton, $\mathsf{R}_e = e_{\mathrm{R}} = \frac{1}{2}(1+\gamma_5)e$, which is a weak-isospin singlet. This completes a specification of the charged weak currents.

To incorporate electromagnetism, Glashow defines a weak hypercharge, Y. Requiring that the Gell-Mann–Nishijima relation for the electric charge, $Q = I_3 + \frac{1}{2}Y$, be satisfied leads to the assignments $Y_{\mathrm{L}} = -1$ and $Y_{\mathrm{R}} = -2$. By construction, the weak-isospin projection I_3 and the weak hypercharge Y are commuting observables.

We now take the (product) group of transformations generated by I and Y to be the gauge group $SU(2)_L \otimes U(1)_Y$ of the theory. To construct the theory, we introduce the gauge fields

$$\begin{array}{cl} b^1_\mu, b^2_\mu, b^3_\mu & \text{for } SU(2)_L, \\ \mathcal{A}_\mu & \text{for } U(1)_Y. \end{array} \qquad 11.$$

Evidently, the Lagrangian for the theory may be written as

$$\mathcal{L} = \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{leptons}}, \qquad 12.$$

where the kinetic term for the gauge fields is

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4}F^l_{\mu\nu}F^{l\mu\nu} - \frac{1}{4}f_{\mu\nu}f^{\mu\nu}, \qquad 13.$$

and the field-strength tensors are $F^l_{\mu\nu} = \partial_\nu b^l_\mu - \partial_\mu b^l_\nu + g\varepsilon_{jkl}b^j_\mu b^k_\nu$ for the $SU(2)_L$ gauge fields and $f_{\mu\nu} = \partial_\nu \mathcal{A}_\mu - \partial_\mu \mathcal{A}_\nu$ for the $U(1)_Y$ gauge field. The matter term is

$$\mathcal{L}_{\text{leptons}} = \bar{\mathsf{R}}_e i\gamma^\mu \left(\partial_\mu + \frac{ig'}{2}\mathcal{A}_\mu Y\right)\mathsf{R}_e + \bar{\mathsf{L}}_e i\gamma^\mu \left(\partial_\mu + \frac{ig'}{2}\mathcal{A}_\mu Y + \frac{ig}{2}\boldsymbol{\tau}\cdot\mathbf{b}_\mu\right)\mathsf{L}_e. \qquad 14.$$

The coupling of the weak-isospin group $SU(2)_L$ is called g, as in the Yang–Mills theory, and the coupling constant for the weak-hypercharge group $U(1)_Y$ is denoted $g'/2$; the factor 1/2 is chosen to simplify later expressions. Similar structures appear for the hadronic weak interactions, now expressed in terms of quarks. The universal strength of charged-current interactions follows from the fact that both the left-handed quarks and the left-handed leptons reside in weak-isospin doublets.

The theory of weak and electromagnetic interactions described by the Lagrangian in Equation 13 is not a satisfactory one, for two immediately obvious reasons. It contains four massless gauge bosons (b^1, b^2, b^3, and $\mathcal{A}$), whereas nature has but one, the photon. In addition, Equation 14 represents a massless electron; it lacks the $-m_e\bar{e}e$ term of the QED Lagrangian of Equation 8, and for good reason. A fermion mass term links left-handed and right-handed components: $\bar{e}e = \frac{1}{2}\bar{e}(1-\gamma_5)e + \frac{1}{2}\bar{e}(1+\gamma_5)e = \bar{e}_R e_L + \bar{e}_L e_R$. The left-handed and right-handed components of the electron transform differently under $SU(2)_L$ and $U(1)_Y$, so an explicit fermion mass term would break the $SU(2)_L \otimes U(1)_Y$ gauge invariance of the theory. Such a mass term is forbidden.

3.3. Insights from Superconductivity

How gauge bosons can acquire mass is a conundrum both for the Yang–Mills theory as a description of nuclear forces and for the $SU(2)_L \otimes U(1)_Y$ theory as a description of the weak and electromagnetic interactions. An important general insight is that the symmetries of the laws of nature need not be manifest in the outcome of those laws. Hidden (or secret) symmetries are all around us in the everyday world—for example, in the ordered structures of crystals and snowflakes or in the spontaneous magnetization of a ferromagnetic substance, configurations that belie the O(3) rotation symmetry of electromagnetism. (See Reference 71 for an interesting tour of spontaneous symmetry breaking in many physical contexts.) The common feature of these phenomena is that the symmetry exhibited by the state of lowest energy, the vacuum, is not the full symmetry of the theory. In addition, the vacuum is degenerate, characterized by many states of the same energy, and the choice of any one is aleatory.

As it happens, superconductivity, a rich and fascinating phenomenon from condensed matter physics, points the way to understanding how gauge bosons can acquire mass. In 1911, shortly after he succeeded in liquefying helium and, therefore, could conduct experiments at unprecedented low temperatures, Kamerlingh Onnes (72) observed the sudden vanishing of electrical resistance in a sample of mercury cooled to 4.2 K. This first miracle of superconductivity is of immense technological importance, not least in the magnets that are essential components of the LHC.

The second miracle, which for me marks superconductivity as truly extraordinary, was discovered in 1933 by Meissner & Ochsenfeld (73; for an English translation, see Reference 74): Magnetic flux is excluded from the superconducting medium. A typical penetration depth (75) is on the order of 10 μm. This means that, within the superconductor, the photon has acquired a mass. Here is the germ of the idea that leads to understanding how the force particles in gauge theories could be massive: QED is a gauge theory, and under the special circumstances of a superconductor, the normally massless photon becomes massive, whereas electric charge remains a conserved quantity.

Two decades would pass before the idea would be fully formed and ready for application to theories of the fundamental interactions. The necessary developments included the elaboration of relativistic quantum field theory and the full realization of QED, a focus on the consequences of spontaneous symmetry breaking, the emergence of informative theories of superconductivity, and attention to the special features of gauge theories.

3.4. Spontaneous Symmetry Breaking

Goldstone (76) obtained a key insight into hidden symmetry in field theory by considering the Lagrangian for two scalar fields ϕ_1 and ϕ_2:

$$\mathcal{L} = \frac{1}{2}[(\partial_\mu \phi_1)(\partial^\mu \phi_1) + (\partial_\mu \phi_2)(\partial^\mu \phi_2)] - V(\phi_1^2 + \phi_2^2). \qquad 15.$$

The Lagrangian is invariant under the group SO(2) of rotations in the $\phi_1 - \phi_2$ plane. It is informative to consider the effective potential

$$V(\boldsymbol{\phi}^2) = \frac{1}{2}\mu^2 \boldsymbol{\phi}^2 + \frac{1}{4}|\lambda|(\boldsymbol{\phi}^2)^2, \qquad 16.$$

where $\phi = \binom{\phi_1}{\phi_2}$ and $\boldsymbol{\phi}^2 = \phi_1^2 + \phi_2^2$, and distinguish two cases.

A positive value of the parameter $\mu^2 > 0$ corresponds to the ordinary case of unbroken symmetry. The unique minimum, corresponding to the vacuum state, occurs at $\langle\phi\rangle_0 = \binom{0}{0}$, so for small oscillations the Lagrangian takes the form

$$\mathcal{L}_{\text{so}} = \frac{1}{2}[(\partial_\mu \phi_1)(\partial^\mu \phi_1) - \mu^2 \phi_1^2] + \frac{1}{2}[(\partial_\mu \phi_2)(\partial^\mu \phi_2) - \mu^2 \phi_2^2], \qquad 17.$$

which describes a pair of scalar particles with common mass μ. Thus, the introduction of a symmetric interaction preserves the spectrum of the free theory with $|\lambda| = 0$.

For the choice $\mu^2 < 0$, a line of minima lies along $\langle\phi^2\rangle_0 = -\mu^2/|\lambda| \equiv v^2$, a continuum of distinct vacuum states that are degenerate in energy. The degeneracy follows from the SO(2) symmetry of Equation 15. Designating one state as the vacuum selects a preferred direction in (ϕ_1, ϕ_2) internal symmetry space and amounts to a spontaneous breakdown of the SO(2) symmetry. Let us select as the physical vacuum state the configuration $\langle\phi\rangle_0 = \binom{v}{0}$, as we may always do with a suitable definition of coordinates. Expanding about the vacuum configuration by defining $\phi' \equiv \phi - \langle\phi\rangle_0 \equiv \binom{\eta}{\zeta}$, we obtain the following Lagrangian for small oscillations:

$$\mathcal{L}_{\text{so}} = \frac{1}{2}[(\partial_\mu \eta)(\partial^\mu \eta) + 2\mu^2 \eta^2] + \frac{1}{2}[(\partial_\mu \zeta)(\partial^\mu \zeta)], \qquad 18.$$

plus an irrelevant constant. There are still two particles in the spectrum. The η particle, associated with radial oscillations, has $(\text{mass})^2 = -2\mu^2 > 0$. The ζ particle, however, is massless. The mass of the η particle may be viewed as a consequence of the restoring force of the potential against radial oscillations. In contrast, the masslessness of the ζ particle is a consequence of the SO(2) invariance of the Lagrangian, which means that there is no restoring force against angular oscillations. It is ironic that the η particle, which here seems so unremarkable, is precisely what emerges as the "Higgs boson" when the hidden symmetry is a gauge symmetry.

The splitting of the spectrum and the appearance of the massless particle are known as the Goldstone phenomenon. Such massless particles, zero-energy excitations that connect possible vacua, are called Nambu–Goldstone bosons. Many occurrences are known in particle, nuclear, and condensed matter physics (77). In any field theory that obeys the usual axioms, including locality, Lorentz invariance, and positive-definite norm on the Hilbert space, if an exact continuous symmetry of the Lagrangian is not a symmetry of the physical vacuum, then the theory must contain a massless spin-0 particle (or particles) whose quantum numbers are those of the broken group generator (or generators) (78).

Goldstone phenomenon: the appearance of massless modes whenever a global continuous symmetry of the Lagrangian is broken, in the sense that the vacuum state does not display the full symmetry of the Lagrangian; one massless scalar or pseudoscalar appears for each broken generator of the full symmetry

This strong statement seemed a powerful impediment to the use of spontaneous symmetry breaking in realistic theories of the fundamental interactions, as the disease of unobserved massless spin-0 particles was added to the disease of massless gauge bosons. Motivated by analogy with the plasmon theory of the free-electron gas, Anderson (79, cf. JS Bach's *Easter Cantata*, BWV 4, versus IV: "Wie ein Tod den andern fraß") put forward a prescient conjecture that one zero-mass ill might cancel the other and make possible a realistic Yang–Mills theory of the strong interactions.

The decisive contributions came at a time of intense interest in superconductivity—that is, in the intricacies of the Bardeen–Cooper–Schrieffer (BCS) theory (80) and in understanding the role of symmetry breaking in the Meissner effect. From the remove of a half-century, it seems to me that preoccupation with the microscopic BCS theory might have complicated the search for a cure for the massless gauge bosons. An easier path is to analyze the phenomenological Ginzburg–Landau (81, 82) description of the superconducting phase transition in the framework of QED. It is then easy to see how the photon acquires mass in a superconducting medium (see, e.g., problem 5.7 of Reference 9, section 21.6 of Reference 83, and the so-called Abelian Higgs model of Reference 84). But that is hindsight and speculation!

Searching for a solution to the problem of massless gauge bosons in field theory, Englert & Brout (18), Higgs (19, 20), and Guralnik et al. (21) showed that gauge theories are different. They do not satisfy the assumptions on which Goldstone's theorem is based, although they are respectable field theories. Recall that to quantize electrodynamics, an exemplary gauge theory, one must choose between the covariant Gupta–Bleuler formalism with its unphysical indefinite-metric states and quantization in a physical gauge for which manifest covariance is lost. Through these authors' work, we understand that the would-be Goldstone bosons that correspond to broken generators of a gauge symmetry become the longitudinal components of the corresponding gauge bosons. What remains as scalar degrees of freedom is an incomplete multiplet—defined by the unbroken generators of the gauge symmetry—of massive particles that we call Higgs bosons.

These authors' collective insight did not, as many had hoped, give rise to a proper description of the strong nuclear force out of Yang–Mills theory. It did, however, set the stage for the development of the electroweak theory and for plausible, if still speculative, unified theories of the strong, weak, and electromagnetic interactions.

It is inaccurate to say that the work of these theorists solved a problem in the Standard Model—the Standard Model did not yet exist! Indeed, they were not concerned with the weak interactions, and the implications for fermion mass shifts are mentioned only in passing. (Recall that for nonchiral theories such as QED and the Yang–Mills theory, the origin of fermion masses does not arise, in the sense that fermion mass is consistent with the gauge symmetry.) Rather, these theorists' work can be said to have triggered the conception of the electroweak theory, which is a very considerable achievement.

Following the discovery of the Higgs boson of the electroweak theory, Englert (85) and Higgs (86) shared the 2014 Nobel Prize for Physics. Guralnik & Hagen (87) have published a memoir of their work. In addition, several of the leading actors in the discovery of spontaneous gauge symmetry breaking as an origin of particle mass have described their personal involvement: Anderson,[4] Englert (88), Guralnik (89, 90), and Higgs (91, 92). Their words carry a special fascination.

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

3.5. The Electroweak Theory and the Standard Model Higgs Boson

In the late 1960s, Weinberg (2) and Salam (3) used the new insights into spontaneous breaking of gauge symmetry to complete the program set out by Glashow (1), described in Section 3.2. The construction of the spontaneously broken $\mathrm{SU(2)_L} \otimes \mathrm{U(1)}_Y$ theory of the weak and electromagnetic interactions is discussed in detail in several papers, including section 2 of Reference 24, so we focus here on a few important conceptual matters.

If $\mathrm{SU(2)_L} \otimes \mathrm{U(1)}_Y$ proves to be the apt choice of gauge symmetry for a theory of weak and electromagnetic interactions, then that symmetry must be hidden, or broken down to the $\mathrm{U(1)_{em}}$ symmetry we observe manifestly. The simple choice made by Weinberg and Salam, which now has significant empirical support, is to introduce a complex weak-isospin doublet of auxiliary scalar fields, and to contrive their self-interactions to create a degenerate vacuum that does not exhibit the full $\mathrm{SU(2)_L} \otimes \mathrm{U(1)}_Y$ symmetry. Before spontaneous symmetry breaking, we count eight degrees of freedom among the four massless gauge bosons and four degrees of freedom for the scalar fields. After spontaneous breaking of $\mathrm{SU(2)_L} \otimes \mathrm{U(1)}_Y \to \mathrm{U(1)_{em}}$, following the path reviewed in Section 3.4, the scalar field obtains the vacuum expectation value $\langle\phi\rangle_0 = \binom{0}{v/\sqrt{2}}$, where $v = (G_F\sqrt{2})^{-\frac{1}{2}} \approx 246$ GeV, to reproduce the low-energy phenomenology.

The charged gauge bosons, $W^\pm$, which mediate the V−A charged-current interaction, acquire mass $gv/2$. The neutral gauge bosons of Equation 11 mix to yield a massive ($M_Z = M_W / \cos\theta_W$) neutral gauge boson, Z^0, that mediates a hitherto unknown weak-neutral-current interaction and a massless photon, γ. (The weak mixing angle θ_W, which parameterizes the mixing of b^3_μ and A_μ, is determined from experiment.) The photon has pure vector couplings, as required, whereas Z^0 has a mix of vector and axial–vector couplings that depend on the quantum numbers of the fermion in question. Eleven of the 12 bosonic degrees of freedom now reside in the vector bosons: 3×3 massive bosons $+1\times2$ massless photon. The last degree of freedom corresponds to the Higgs boson: It is a massive scalar, but the Weinberg–Salam theory does not fix its mass. Because the Higgs boson and the longitudinal components of the gauge bosons share a common origin, the Higgs boson plays an essential role in ensuring a sensible high-energy behavior of the electroweak theory (93). All this is fixed by the construction of the theory: Once the representation of the auxiliary scalar fields is chosen and the weak mixing parameter determined, all the couplings of gauge bosons to fermions and couplings among gauge and Higgs bosons are set. The new neutral-current interactions among the leptons are flavor diagonal. Also, to this point, we have solved only one of the outstanding problems of the unbroken $\mathrm{SU(2)_L} \otimes \mathrm{U(1)}_Y$ theory: the masslessness of all the gauge bosons. What of the fermions? Weinberg and Salam saw the possibility to generate fermion masses in the spontaneously broken theory by adding to the Lagrangian a gauge-invariant interaction, $\mathcal{L}_{\mathrm{Yukawa}} = -\zeta_e[\bar{\mathsf{R}}_e(\phi^\dagger\mathsf{L}_e) + (\bar{\mathsf{L}}_e\phi)\mathsf{R}_e]$, where the Yukawa coupling, ζ_e, is a phenomenological parameter. When

[4] See the interview with Philip Anderson by Alexei Kojevnikov on November 23, 1999, Niels Bohr Library and Archives, American Institute of Physics (**http://j.mp/1DYYfsM**), in particular the passage beginning "Now, during that year in Cambridge . . .."

the gauge symmetry is hidden, the Yukawa term becomes

$$\mathcal{L}_{\text{Yukawa}} = -\frac{\zeta_e v}{\sqrt{2}}\bar{e}e - \frac{\zeta_e H}{\sqrt{2}}\bar{e}e, \tag{19}$$

where H is the Higgs boson. The electron has acquired a mass $\zeta_e v/\sqrt{2}$, and the $He\bar{e}$ coupling is $-im_e/v$. It is pleasing that the electron mass arises spontaneously, but frustrating that the parameter ζ_e must be put in by hand and does not emerge from the theory. The same strategy carries over for all the quarks and charged leptons, and may also be considered the origin of the parameters of the quark-mixing matrix.

If the Higgs field is the source of the quark and charged-lepton masses, that does not mean that the Higgs boson is the source of all mass in the Universe, as is frequently stated—even by physicists. The overwhelming majority of the visible mass in the Universe is in the form of atoms, and most of that is made up of nucleon mass, which arises as confinement energy in QCD (94). Electroweak symmetry breaking is decidedly a minor player.

Only three quark flavors (u, d, and s) were known when the electroweak theory was formulated. The weak-isospin doublet $\binom{u}{d\cos\theta_C + s\sin\theta_C}_L$, where θ_C is the Cabibbo angle, captured the known structure of the hadronic charged-current interaction and expressed the universal strength of quark and lepton interactions. Within the Weinberg–Salam framework, however, this single quark doublet gives rise to flavor-changing $s \leftrightarrow d$ neutral-current interactions that are not observed in nature. Glashow, Iliopoulos, and Maiani (GIM) (95) noted that the unwanted interactions could be cancelled by introducing a second quark doublet, $\binom{c}{s\cos\theta_C - d\sin\theta_C}_L$, involving a new "charmed" quark and the orthogonal combination of d and s. The absence of flavor-changing neutral currents generalizes to more (complete) quark doublets, and is a striking feature of the experimental data.

To keep expressions compact, I outline here a theory of a single generation of leptons; the other lepton families are included as simple copies. However, a theory of leptons alone would be inconsistent. In our left-handed world, each doublet of leptons must be accompanied by a color-triplet weak-isospin doublet of quarks, in order that the theory be anomaly free—in other words, that quantum corrections respect the symmetries on which the theory is grounded (96).

Since its invention, the electroweak theory has been supported again and again by new observations, in many cases arising from experiments conceived or reoriented explicitly to test the electroweak theory. I treated this question in some detail in section 3 of Reference 24, to which I refer the reader for specific references. It suffices here to mention some of the major supporting elements. The first great triumph of the electroweak theory was the discovery of weak neutral currents. This discovery was soon followed by the discovery of charm (hidden first, then open), which was required in the framework of the electroweak theory, once neutral currents had been observed. The discovery of the W and Z bosons was the second great triumph of the electroweak theory. Experiments also brought new evidence of richness, including the discovery of the τ lepton as well as evidence for a distinct τ neutrino and the discovery of the b quark. The t quark completed a third quark generation; the t quark mass became an essential input to quantum corrections to predictions for precisely measured observables. Moreover, finding a third quark generation opened the way to understanding, at least at an operational level, the systematics of *CP* violation. Highly detailed studies at many laboratories confirmed the predictions of the electroweak theory to an extraordinary degree.

As the electroweak theory emerged as a new law of nature, the question of how the electroweak symmetry was hidden became central. Although the default option—the one emphasized in textbooks—was an elementary scalar Higgs boson, electroweak symmetry breaking, through some sort of new strong dynamics, or as a message from extra spatial dimensions, or as an emergent

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

GIM mechanism: observation by Glashow, Iliopoulos, and Maiani (95) that flavor-changing neutral-current interactions vanish at tree level and are strongly inhibited at higher orders, provided that quarks (and leptons) occur in $SU(2)_L$ doublets; argues for the necessity of the charm quark

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 phenomenon arising from strong interactions among the weak bosons, received attention from both theory and experiment.

4. AFTER THE DISCOVERY: OUTLOOK

The most succinct summary we can give is that the data from the ATLAS and CMS experiments are developing as if electroweak symmetry is broken spontaneously through the work of elementary scalars, and that the emblem of that mechanism is the Standard Model Higgs boson. I refer to References 22 and 23 for details and to Reference 17 for perspective.

The bare facts are these: The LHC experiments have found a new, unstable particle H, with a mass of approximately 125 GeV. It decays into $\gamma\gamma$, W^+W^-, and Z^0Z^0 in approximately the proportions expected for a Standard Model Higgs boson. The new particle is narrow for its mass, with the current bounds measured in tens of MeV. The dominant production mechanism has characteristics that are compatible with gluon fusion through a heavy-quark loop, as foreseen. Topological selections have identified a subsidiary mechanism compatible with vector-boson fusion. Some evidence has been presented for the decays $H \to b\bar{b}$ and $\tau^+\tau^-$. No decays that entail lepton flavor violation have been observed. The new particle does not have spin 1; studies of decay

ANSWERS TO SOME QUESTIONS POSED BEFORE THE LHC EXPERIMENTS [a]

1. Q: What is the agent that hides the electroweak symmetry? Specifically, is there a Higgs boson? Might there be several?
A: To the best of our knowledge, $H(125)$ displays the characteristics of a Standard Model Higgs boson, an elementary scalar. Searches will continue for other particles that may play a role in electroweak symmetry breaking.

2. Q: Is the "Higgs boson" elementary or composite? How does the Higgs boson interact with itself? What triggers electroweak symmetry breaking?
A: We have not yet found any evidence that $H(125)$ is other than an elementary scalar. Searches for a composite component will continue. The Higgs boson self-interaction is almost certainly out of the reach of the LHC; it is a very challenging target for future, very high energy accelerators. We do not yet know what triggers electroweak symmetry breaking.

3. Q: 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?
A: The experimental evidence suggests that $H(125)$ couples to $t\bar{t}$, $b\bar{b}$, and $\tau^+\tau^-$, so the answer is probably yes. All these are third-generation fermions, so even if the evidence for these couplings becomes increasingly robust, we will want to see evidence that H couples to lighter fermions. The most likely candidate, perhaps in high-luminosity LHC running, is for the $H\mu\mu$ coupling, which would already show that the third generation is not unique in its relation to H. Ultimately, to show that spontaneous symmetry breaking accounts for electron mass, and thus enables compact atoms, we will want to establish the $He\bar{e}$ coupling. Doing so will be extraordinarily challenging because of the minute branching fraction.

10. Q: What lessons does electroweak symmetry breaking hold for unified theories of the strong, weak, and electromagnetic interactions?
A: Establishing that scalar fields drive electroweak symmetry breaking will encourage the already standard practice of using auxiliary scalars to hide the symmetries that underlie unified theories.

[a] The question numbers correspond to those in Reference 24.

angular distributions and correlations among decay products strongly favor spin parity 0^+ over 0^-; and whereas spin 2 has not been excluded in the most general case, that assignment is implausible.

As one measure of the progress the discovery of the Higgs boson represents, let us consider some of the questions I posed before the LHC experiments (see the sidebar). To close, I offer a revised list of questions (see Future Issues) to build on what our first look at the Higgs boson has taught us. In the realms of refined measurements, searches, and theoretical analyses and imagination, great opportunities lie before us!

SUMMARY POINTS

1. The ATLAS and CMS Collaborations, working at CERN's LHC, have discovered a new particle, $H(125\text{ GeV})$, that matches the profile of the Higgs boson of the electroweak theory.
2. Observation of decays into the weak bosons, $H \to W^+W^-$ and $H \to Z^0Z^0$, establishes a role for the Higgs boson in hiding the electroweak symmetry and endowing the weak bosons with mass.
3. Evidence for the decays $H \to b\bar{b}$ and $H \to \tau^+\tau^-$, together with characteristics of $H(125)$ production that implicate gluon fusion through a t quark loop, suggest that the Higgs boson also plays a role in giving mass to the fermions.
4. It will be important to show that the new particle couples to quarks and leptons of the first two generations and to test its role in generating their masses.
5. If the electron mass, in particular, does arise from the vacuum expectation value of the Higgs field, we will have a new understanding of why compact atoms exist, why valence bonding is possible, and why liquids and solids can form (see section 4.4.2 of Reference 24).
6. The spin parity of the new particle, which is strongly indicated as 0^+, favors the interpretation as an elementary scalar.
7. Even after its apparent completion by the observation of a light Higgs boson, the electroweak theory raises questions. An outstanding issue is why the electroweak scale is so much smaller than other plausible physical scales, such as the unification scale and the Planck scale.
8. It is possible that the Higgs boson experiences new forces or decays into hitherto unknown particles.

FUTURE ISSUES

1. How closely does $H(125)$ hew to the expectations for a Standard Model Higgs boson? Does it have any partners that contribute appreciably to electroweak symmetry breaking?
2. Do the HZZ and HWW couplings indicate that $H(125)$ is solely responsible for electroweak symmetry breaking, or is it only part of the story?

3. Does the Higgs field give mass to fermions beyond the third generation? Does $H(125)$ account quantitatively for the quark and lepton masses? What sets the masses and mixings of the quarks and leptons?
4. What stabilizes the Higgs boson mass below 1 TeV?
5. Does the Higgs boson decay to new particles, or via new forces?
6. What will be the next symmetry recognized in nature? Is nature supersymmetric? Is the electroweak theory part of some larger edifice?
7. Are all the production mechanisms as expected?
8. Is there any role for strong dynamics? Is electroweak symmetry breaking related to gravity through extra space-time dimensions?
9. What lessons does electroweak symmetry breaking hold for unified theories of the strong, weak, and electromagnetic interactions?
10. What implications does the value of the $H(125)$ mass have for speculations that go beyond the Standard Model, and for the range of applicability of the electroweak theory?

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