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SUPERCOLLIDER PHYSICS

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Abstract

Gauge theories of the strong, weak, and electromagnetic interactions organize our knowledge of particle physics and achieve a wide-ranging synthesis of physical phenomena, but they raise some fundamental questions that we cannot yet answer. Important clues to a more complete understanding are to be found on the energy scale of 1 TeV, where the kinship between the weak and electromagnetic interactions and the origin of mass will be illuminated. Exploration of the 1-TeV scale is an essential goal for particle physics.

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1 Introduction

About a decade ago, I gave a series of talks on Supercollider Physics[1)] in which I expressed a desire to return in ten years to discuss the results of experiments to explore the 1-TeV scale. The organizers of this Rencontre de la Vallée d'Aoste have granted my wish. Unfortunately, the way toward an understanding of electroweak symmetry breaking has been steeper than we hoped. We do not have the experimental results from the SSC that I dreamed of, and we will not have them soon. Our other hope, the Large Hadron Collider at CERN, is still about ten years away.

The demise of the SSC and the delay in exploring the 1-TeV scale do not alter the fact that the problem of electroweak symmetry breaking—of the origin of mass—is a defining issue of particle physics today. I will leave it to Peter Jenni[2)] to distill ten years of effort on what experiments can be done and what heroic efforts they will require. I would like to reflect more generally on why physics at the 1-TeV scale is of such compelling interest.

Our focus on the problem of mass has been influenced by twenty-five years of important progress in particle physics. Indeed, at this conference, I have been struck by the fact that nearly all the topics discussed here were literally unknown a quarter of a century ago, or have been entirely transfigured by the developments of our recent history. That is an impressive observation about the vitality of particle physics, and we should keep that intellectual vitality in the foreground in these times when finding the means to do our science is a challenge.

I want to introduce the problem of mass—the problem of electroweak symmetry breaking—through an allegory.

2 An Allegory of Three Worlds

The Crystal World. Think of what it would be like to be a very tiny physicist, a millionth of a millionth of normal size, living and working in the recesses of a magnetic crystal of iron. The laws of electromagnetism that, together with quantum mechanics, determine the shape of the crystal world, display an exact rotational symmetry—an exact $O(3)$ symmetry. But this picophysicist would have a hard time learning that nature favors no direction over any other, because on every street corner of the picoworld there stands a compass, and the needle of every compass points the same way. The compasses are monuments to the way things were at a time that no one remembers, just after this world came out of the fiery furnace.

All these compass needles make an unchanging magnetic field of about 2 teslas that pervades the whole world, surrounding the picophysicist and his instruments. If the picophysicist's picoinstruments were not much affected by the magnetic field of the regimented compass needles, our tiny colleague might find that the laws of nature look approximately the same from every direction. But if the picoapparatus were strongly influenced by magnetism, the picophysicist might miss the idea of rotational symmetry altogether. In any event, picoexperiments would never reveal that rotation symmetry is exact.

Nor could the picophysicist show directly that the orientation of the compasses, the emblem of stability and order in this world, is determined by happenstance. (You can imagine the cultural significance that would have accumulated over the years for these compass needles all pointing in the same direction.) Too small by far to reorient all the compass needles at the same time, the picophysicist would find it an impossible task to show that any other uniform alignment is equivalent. The picoscholar could at best conjecture the idea that there is a symmetry hidden behind the order of his world and test the idea through its consequences.

At some considerable peril to our miniature colleague, we macroscopic outsiders could be of assistance. We could raise the crystal's temperature to 1040 kelvins, until heat's random motion disordered the needles. Then, as the crystal cooled, we could watch all the street-corner compasses settle into a new pattern, equally regimented but differently aligned. The preferred direction would change on every cycle of heating and cooling, with no memory of what had gone before. It would reveal by its aleatory nature the hidden rotational symmetry that governs the magnetic crystal picoworld.

A Perfect World. This is a world of divine perfection, where time and space flow like the still waters of a deep river. No street corners are marked by prehistoric monuments, no direction markers point the one true way. Vast expanses of space and time are unrelieved flatness.

In this egalitarian world, matter particles and force carriers all dash about at the speed of light, exchanging information in brief encounters, never stopping to form lasting associations. Liaisons are here today, gone today. There are no atoms, no complex structures, no condensed-matter physicists. (I think of it as Leon Lederman's Paradise.) All particles are brothers and all forces are one. It is a world of perfect symmetry and complete disorder—not to say anarchy, for symmetry rules with so heavy a hand that it imposes an unrelenting sameness, a stability in mutability. Everything is interchangeable. It is a perfectly boring world.

The Third World. In a world of diversity, space is not punctuated by a crystal framework, but runs, like time, with unbroken continuity. In the eyes of nature's laws, no time and place is preferred to any other; no direction is superior to the rest. Yet this is a world of distinctions, a world in which differences matter.

Quarks stand apart from leptons. Leptons are free, quarks confined. Every quark, every lepton has a distinct personality, every force a peculiar character. Some changes—the actions of forces—are everyday events. Others happen once in a proton's lifetime.

Some matter particles and force carriers have weight, can come to rest. Others, weightless, are in perpetual motion at the speed of light. Some particles that weigh are ephemeral; they transmute, decay. Other particles can live forever. In this world of mass, composites form, stable structures are commonplace. Accumulations of matter ripple the fabric of spacetime.

Like the magnetic crystal, this is a world of bias. A pervasive tilt, set at random when the world cooled from a state of symmetry, disorder, and perfection, veils an exact symmetry. It distinguishes up quarks from down, electrons from neutrinos; it invests particles with mass.

This world is ours. Like our picocolleague in the crystal world, we cannot hope to undo the bias or change it for another. We must peer through the veil to discover the source of bias, to learn what hides the symmetries that lie behind the order.

In the magnetic crystal, the state of lowest energy—the vacuum state—does not display the full rotational invariance of electromagnetism. Below the Curie temperature, the magnetic interaction among the tiny dipoles overcomes thermal agitation and causes an alignment, the selection of a preferred axis. Only rotations about the axis of magnetization leave the crystal unchanged in appearance. The full $O(3)$ rotation invariance is hidden, or spontaneously broken. That reduced rotational invariance of the crystal world corresponds to the spontaneously broken $SU(2)_L \otimes U(1)_Y$ gauge symmetry of our world. The spontaneous magnetization in the crystal world corresponds in our world to a reference in weak isospin space. Heisenberg's Hamiltonian for the ferromagnet corresponds to our quantum field theory, including the Higgs potential.

If the perfect world is a hot version of our own world, then what happens at very high energies—what happens in that state of perfection at high temperatures—is encoded in the fine structure constant, the number that determines the strength of electromagnetism. Because if at very high energies all forces are equal and have the same strength, and the residual differences we see in our low-energy world are due to the hiding of the symmetry—the perfection—that is evident at high temperatures, then everything that happens between here and there, everything that happens from very high temperatures down to low temperatures influences the value of this number that determines the size of atoms and the strength of electromagnetism. Top, bottom and even the supersymmetric partners that we might some day discover determine the dimensions and character of this world that we live in. Unified theories such as the one I've speculated about just now[3] are not empty, untestable exercises in metaphysics. They have consequences for our world.

3 Hiding a Gauge Symmetry

Although the parallels between our world and the ferromagnetic crystal are strong, the most apt analogy for the hiding of the electroweak gauge symmetry is found in superconductivity. In the Ginzburg-Landau description[4] of the superconducting phase transition, a superconducting material is regarded as a collection of two kinds of charge carriers: normal, resistive carriers, and superconducting, resistanceless carriers.

In the absence of a magnetic field, the free energy of the superconductor is related to the free energy in the normal state through

$$G_{\mathrm{super}}(0) = G_{\mathrm{normal}}(0) + \alpha\, |\psi|^2 + \beta\, |\phi|^4 \quad , \tag{3.1}$$

where α and β are phenomenological parameters and $|\psi|^2$ is an order parameter corresponding to the density of superconducting charge carriers. The parameter β is non-negative, so that the free energy is bounded from below.

Above the critical temperature for the onset of superconductivity, the parameter α is positive and the free energy of the substance is supposed to be an increasing function of the density of superconducting carriers, as shown in Figure 1(a). The state of minimum energy, the vacuum state, then corresponds to a purely resistive flow, with no superconducting carriers active. Below the critical temperature, α is negative and the free energy is minimized when $\psi = \psi_0 \neq 0$, as illustrated in Figure 1(b).

This is a nice cartoon description of the superconducting phase transition, but there is more. In an applied magnetic field $\vec{H}$, the free energy is

$$G_{\mathrm{super}}(\vec{H}) = G_{\mathrm{super}}(0) + \frac{\vec{H}^2}{8\pi} + \frac{1}{2m^\star}|-i\hbar\nabla\psi - (e^\star/c)\vec{A}\psi|^2 \quad , \tag{3.2}$$

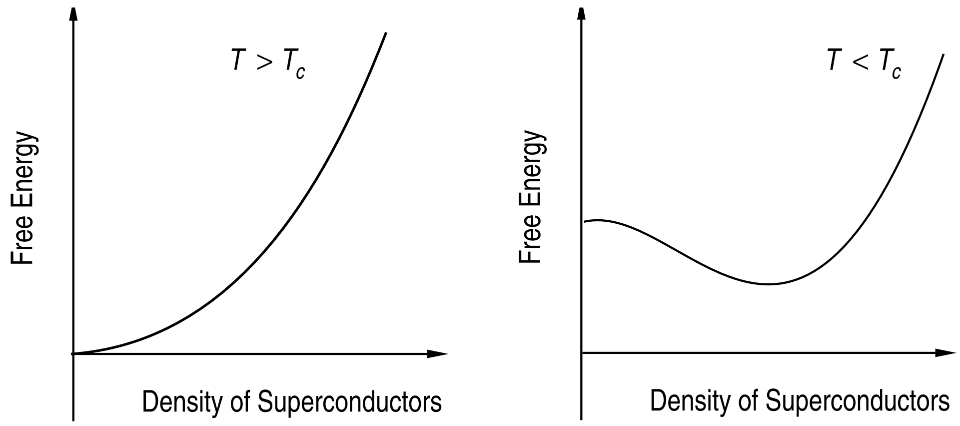


Figure 1: Ginzburg-Landau description of the superconducting phase transition.

where $e^\star$ and $m^\star$ are the charge (-2 units) and effective mass of the superconducting carriers. In a weak, slowly varying field $\vec{H} \approx 0$, when we can approximate $\psi \approx \psi_0$ and $\nabla\psi \approx 0$, the usual variational analysis leads to the equation of motion,

$$\nabla^2 \vec{A} - \frac{4\pi e^\star}{m^\star c^2} \left|\psi_0\right|^2 \vec{A} = 0 \quad , \tag{3.3}$$

the wave equation of a massive photon. In other words, the photon acquires a mass within the superconductor. This is the origin of the Meissner effect, the exclusion of a magnetic field from a superconductor. More to the point, for our purposes, it shows how a symmetry-hiding phase transition can lead to a massive gauge boson.

To give masses to the intermediate bosons of the weak interaction, we take advantage of a relativistic generalization of the Ginzburg-Landau phase transition known as the Higgs mechanism.[5)] We introduce auxiliary scalar fields, with gauge-invariant interactions among themselves and with the fermions and bosons of the electroweak theory. We then arrange their self-interactions so that the vacuum state corresponds to a broken-symmetry solution. As a result, the W and Z bosons acquire masses, as auxiliary scalars assume the role of the third (longitudinal) degrees of freedom of what had been massless gauge bosons. The quarks and leptons acquire masses as well, from their Yukawa interactions with the scalars. Finally, there remains as a vestige of the spontaneous breaking of the symmetry a massive, spin-zero particle, the Higgs boson. 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.

It is remarkable that the resulting theory has been tested at distances ranging from about 10^{-16} cm to about 4×10^{20} cm, especially when we consider that classical electrodynamics has its roots in the tabletop experiments that gave us Coulomb's law. These basic ideas were modified in response to the quantum effects observed in atomic experiments. High-energy physics experiments both inspired and tested the unification of weak and electromagnetic interactions. At distances longer than common experience, electrodynamics—in the form of the statement that the photon is massless—has been tested in measurements of the magnetic fields of the planets. With additional assumptions, the observed stability of the Magellanic clouds provides evidence that the photon is massless over distances of about 10^{22} cm. The wonderful agreement between the electroweak theory and experiments at the Z^0-pole has been summarized here at La Thuile by Michael Koratzinos[6)] and Riccardo Barbieri[7)]. We eagerly await the new tests that will become possible when the top quark is discovered and the W-boson mass is measured to greater precision at Fermilab and LEP200.

4 Why a Higgs Boson Must Exist

How can we be sure 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. An illuminating example is provided by the reaction

$$e^+e^- \rightarrow W^+W^-, \tag{4.1}$$

which is described in lowest order by the four Feynman graphs in Figure 2. 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). However, 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 $s^{1/2}$ for the production of longitudinally polarized gauge bosons. The resulting divergence 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 "wrong-helicity" amplitudes are always proportional to 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

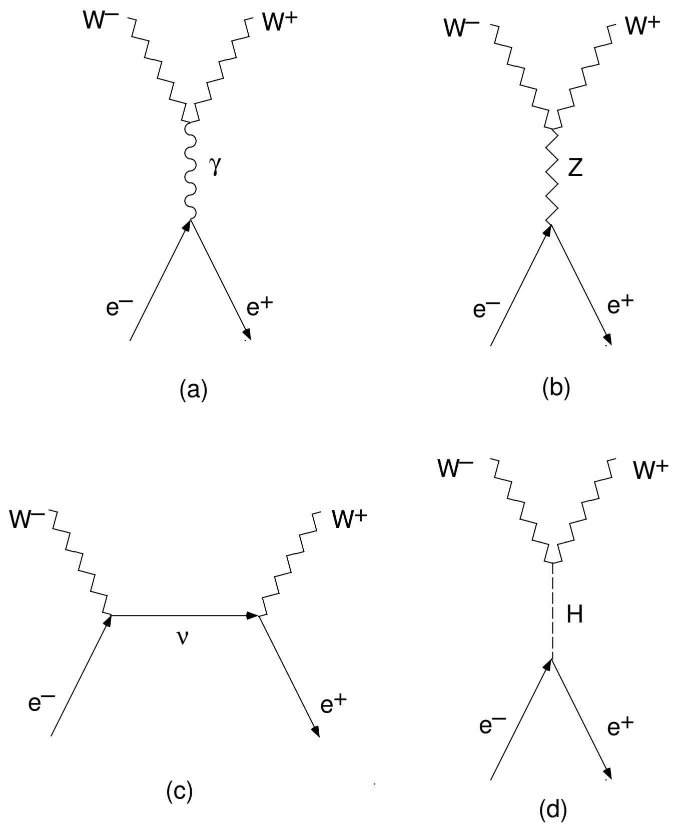


Figure 2: Lowest-order contributions to the e^+e^- scattering amplitude.

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.

5 Strong Scattering of Gauge Bosons

The threshold behavior of the partial-wave amplitudes for gauge-boson scattering follows generally from chiral symmetry.[8)] The partial-wave amplitudes a_{IJ} of definite isospin I and angular momentum J are given by

$$\begin{aligned} a_{00} &\sim \quad G_F s/8\pi\sqrt{2} \quad \text{attractive,} \\ a_{11} &\sim \quad G_F s/48\pi\sqrt{2} \quad \text{attractive,} \\ a_{20} &\sim -G_F s/16\pi\sqrt{2} \quad \text{repulsive.} \end{aligned} \tag{5.1}$$

Unless the mass M_H of the Higgs boson is less than about 1 TeV/c^2, these amplitudes grow to exceed the unitarity bound $|a_{IJ}| < 1$ for $s \approx 2$ GeV^2. This means that the features of strong interactions at GeV energies would characterize electroweak gauge boson interactions at TeV energies.

Consider next the limit of very high energies, for $s \gg M_H^2,\ M_W^2$. 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. Four channels are interesting:

$$W_L^+ W_L^-,\quad Z_L^0 Z_L^0/\sqrt{2},\quad HH/\sqrt{2},\quad HZ_L^0\ , \tag{5.2}$$

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 tree-level amplitudes are all asymptotically constant (*i.e.*, well behaved) and proportional to $G_F M_H^2$. In the high-energy limit,

$$\lim_{s \gg M_H^2} (a_0) \to \frac{-G_F M_H^2}{4\pi\sqrt{2}} \cdot \begin{bmatrix} 1 & 1/\sqrt{8} & 1/\sqrt{8} & 0 \\ 1/\sqrt{8} & 3/4 & 1/4 & 0 \\ 1/\sqrt{8} & 1/4 & 3/4 & 0 \\ 0 & 0 & 0 & 1/2 \end{bmatrix} . \tag{5.3}$$

The matrix has eigenvalues $3/2, 1/2, 1/2, 1/2$. Requiring that the largest eigenvalue respect the partial-wave unitarity condition $|a_0| \leq 1$ yields[9)]

$$M_H < \left(\frac{8\pi\sqrt{2}}{3G_F}\right)^{1/2} \approx 1\ \mathrm{TeV}/c^2 \tag{5.4}$$

as a condition for perturbative unitarity. If the bound is respected, weak interactions remain weak at all energies, and perturbation theory is everywhere reliable.

Both the near-threshold and asymptotic analyses show that if the Higgs-boson mass exceeds 1 TeV/c^2, perturbation theory breaks down, and weak interactions among $W^\pm$, Z, and H become strong on the 1-TeV scale. We interpret this to mean that new phenomena are to be found in the electroweak interactions at energies not much larger than 1 TeV.

6 Parameters in the Higgs Sector

Although the Higgs mechanism shows how masses may be given to the quarks and leptons, as well as the electroweak gauge bosons, the electroweak theory offers no particular insight into the pattern of fermion masses and mixing angles.

Of the nineteen arbitrary parameters of the $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ gauge theory of the fundamental interactions, 3 coupling parameters: α_s, $\alpha_{\rm EM}$, $\sin^2\theta_W$

6 quark masses: m_u, m_d, m_c, m_s, m_t, m_b

3 Cabibbo-Kobayashi-Maskawa angles: θ_1, θ_2, θ_3

1 CP-violating phase: δ

2 parameters of the Higgs potential: μ^2, $|\lambda|$

3 charged-lepton masses: m_e, m_μ, m_τ

1 vacuum phase: Θ,

fully fifteen are associated with the Higgs sector. The situation is not improved by the unification of strong, weak, and electromagnetic interactions. Unification imposes constraints among some parameters, but new parameters arise to describe the spontaneous breakdown of the unifying group into $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$. If neutrinos have mass, still more parameters will be associated with the Higgs sector.

7 Why is the Electroweak Scale Small?

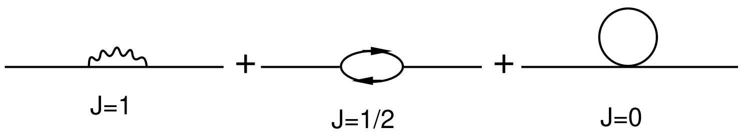
The $SU(2)_L \otimes U(1)_Y$ electroweak theory does not explain how the scale of electroweak symmetry breaking is maintained in the presence of quantum corrections. The problem of the scalar sector can be summarized neatly as follows.[10] The Higgs potential is

$$V(\phi^\dagger\phi) = \mu^2(\phi^\dagger\phi) + |\lambda|\,(\phi^\dagger\phi)^2\ . \tag{7.1}$$

With μ^2 chosen to be less than zero, the electroweak symmetry is spontaneously broken down to the $U(1)$ of electromagnetism, as the scalar field acquires a vacuum expectation value that is fixed by the low-energy phenomenology,

$$\langle\phi\rangle_0 = \sqrt{-\mu^2/2|\lambda|} \equiv (G_F\sqrt{8})^{-1/2} \approx 175\ \text{GeV}\ . \tag{7.2}$$

Beyond the classical approximation, scalar mass parameters receive quantum corrections from loops that contain particles of spins $J = 1, 1/2$, and 0:

$m^2(p^2) = m_0^2 +$ 

(7.3)

The loop integrals are potentially divergent. Symbolically, we may summarize the content of Eq. (7.3) as

$$m^2(p^2) = m^2(\Lambda^2) + Cg^2\int_{p^2}^{\Lambda^2} dk^2 + \cdots\ , \tag{7.4}$$

where Λ defines a reference scale at which the value of m^2 is known, g is the coupling constant of the theory, and the coefficient C is calculable in any particular theory. Instead of dealing with the relationship between observables and parameters of the Lagrangian, we choose to describe the variation of an observable with the momentum scale. In order for the mass shifts induced by radiative 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 cut off the integral.

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 the natural reference scale is the Planck mass,

$$\Lambda \sim M_{\text{Planck}} \approx 10^{19} \text{ GeV} . \tag{7.5}$$

In a unified theory of the strong, weak, and electromagnetic interactions, the natural scale is the unification scale,

$$\Lambda \sim M_U \approx 10^{15}\text{-}10^{16} \text{ GeV} . \tag{7.6}$$

Both estimates are very large compared to the scale of electroweak symmetry breaking (7.2). We are therefore assured that new physics must intervene at an energy of approximately 1 TeV, in order that the shifts in m^2 not be much larger than (7.2).

Only a few distinct scenarios for controlling the contribution of the integral in (7.4) can be envisaged. The supersymmetric solution[7,11,12)] is especially elegant. Exploiting the fact that fermion loops contribute with an overall minus sign (because of Fermi statistics), supersymmetry balances the contributions of fermion and boson loops. In the limit of unbroken supersymmetry, in which the masses of bosons are degenerate with those of their fermion counterparts, the cancellation is exact:

$$\sum_{i=\substack{\text{fermions}\\+\text{bosons}}} C_i \int dk^2 = 0 . \tag{7.7}$$

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$/c^2$.

A second solution to the problem of the enormous range of integration in (7.4) is offered by theories of dynamical symmetry breaking such as technicolor.[13] In technicolor models, the Higgs boson is composite, and new physics arises on the scale of its binding, $\Lambda_{TC} \simeq O(1\text{ TeV})$. Thus the effective range of integration is cut off, and mass shifts are under control.

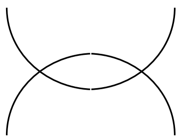
A third possibility is that the gauge sector becomes strongly interacting.[14] This would give rise to WW resonances, multiple production of gauge bosons, and other new phenomena at energies of 1 TeV or so. It is likely that a scalar bound state—a quasi-Higgs boson—would emerge with a mass less than about 1 TeV/c^2.

We cannot avoid the conclusion that some new physics must occur on the 1-TeV scale.

8 Triviality of Scalar Field Theory

The electroweak theory itself provides another reason to expect that discoveries will not end with the Higgs boson. Scalar field theories make sense on all energy scales only if they are noninteracting, or "trivial."[15] The vacuum of quantum field theory is a dielectric medium that screens charge. Accordingly, the effective charge is a function of the distance or, equivalently, of the energy scale. This is the famous phenomenon of the running coupling constant.

In $\lambda\phi^4$ theory (compare the interaction term in the Higgs potential), it is easy to calculate the variation of the coupling constant λ in perturbation theory by summing bubble graphs like this one:



. (8.1)

The coupling constant $\lambda(\mu)$ on a physical scale μ is related to the coupling constant on a higher scale Λ by

$$\frac{1}{\lambda(\mu)} = \frac{1}{\lambda(\Lambda)} + \frac{3}{2\pi^2}\log\left(\Lambda/\mu\right) \ . \tag{8.2}$$

This perturbation-theory result is reliable only when λ is small, but lattice field theory allows us to treat the strong-coupling regime.

In order for the Higgs potential to be stable (*i.e.*, for the energy of the vacuum state not to race off to $-\infty$), $\lambda(\Lambda)$ must not be negative. Therefore we can rewrite (8.2) as an inequality,

$$\frac{1}{\lambda(\mu)} \geq \frac{3}{2\pi^2}\log\left(\Lambda/\mu\right) \ . \tag{8.3}$$

This gives us an *upper bound*,

$$\lambda(\mu) \le 2\pi^2/3\log(\Lambda/\mu) \ , \tag{8.4}$$

on the coupling strength at the physical scale μ. If we require the theory to make sense to arbitrarily high energies—or short distances—then we must take the limit $\Lambda \to \infty$ while holding μ fixed at some reasonable physical scale. In this limit, the bound (8.4) forces $\lambda(\mu)$ to zero. The scalar field theory has become free field theory; in theorist's jargon, it is trivial.

We can rewrite the inequality (8.4) as a bound on the Higgs-boson mass. Rearranging and exponentiating both sides gives the condition

$$\Lambda \le \mu \exp\left(\frac{2\pi^2}{3\lambda(\mu)}\right) \ . \tag{8.5}$$

Choosing the physical scale as $\mu = M_H$, and remembering that, before quantum corrections,

$$M_H^2 = 2\lambda(M_H)v^2 \ , \tag{8.6}$$

where $v = (G_F\sqrt{2})^{-1/2} \approx 246$ GeV is the vacuum expectation value of the Higgs field times $\sqrt{2}$, we find that

$$\Lambda \le M_H \exp\left(\frac{4\pi^2 v^2}{3M_H^2}\right) \ . \tag{8.7}$$

For any given Higgs-boson mass, there is a maximum energy scale $\Lambda^\star$ at which the theory ceases to make sense. The description of the Higgs boson as an elementary scalar is at best an effective theory, valid over a finite range of energies.

If the Higgs boson is relatively light—which would itself require explanation—then the theory can be self-consistent up to very high energies. If the electroweak theory is to make sense all the way up to a unification scale $\Lambda^\star = 10^{16}$ GeV, then the Higgs boson must weigh less than 170 GeV/c^2.[16)]

This perturbative analysis breaks down when the Higgs-boson mass approaches 1 TeV/c^2 and the interactions become strong. Lattice analyses[17)] indicate that, for the theory to make sense up to a few TeV, the mass of the Higgs boson can be no more than about 800 GeV/c^2. Another way of putting this result is that, if the elementary Higgs boson takes on the largest mass allowed by perturbative unitarity arguments, the electroweak theory will be living on the brink of instability.

9 Technicolor

I now consider one example of several possible extensions to the electroweak theory: the technicolor scenario for dynamical symmetry breaking. I select this in part because the other leading candidate, supersymmetry, is so well known, and in part because I find its claim on our attention very powerful.[18)]

The dynamical-symmetry-breaking approach realized in technicolor theories is modeled upon our understanding of the superconducting phase transition. The macroscopic order parameter of the Ginzburg-Landau phenomenology corresponds to the wave function of superconducting charge carriers. As we have seen in §3, it acquires a nonzero vacuum expectation value in the superconducting state. The microscopic Bardeen-Cooper-Schrieffer theory[19)] identifies the dynamical origin of the order parameter with the formation of bound states of elementary fermions, the Cooper pairs of electrons. The basic idea of technicolor is to replace the elementary Higgs boson with a fermion-antifermion bound state. By analogy with the superconducting phase transition, the dynamics of the fundamental technicolor gauge interactions among technifermions generate scalar bound states, and these play the role of the Higgs fields.

The elementary fermions—electrons—and the gauge interactions—QED—needed to generate the scalar bound states are already present in the case of superconductivity. Could a scheme of similar economy account for the transition that hides the electroweak symmetry? Consider an $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ theory of massless up and down quarks. Because the strong interaction is strong, and the electroweak interaction is feeble, we may treat the $SU(2)_L \otimes U(1)_Y$ interaction as a perturbation. For vanishing quark masses, QCD has an exact $SU(2)_L \otimes SU(2)_R$ chiral symmetry. At an energy scale $\sim \Lambda_{QCD}$, the strong interactions become strong, fermion condensates appear, and the chiral symmetry is spontaneously broken to the familiar flavor symmetry:

$$SU(2)_L \otimes SU(2)_R \to SU(2)_V \quad . \tag{9.1}$$

Three Goldstone bosons appear, one for each broken generator of the original chiral invariance. These were identified by Nambu[20)] as three massless pions.

The broken generators are three axial currents whose couplings to pions are measured by the pion decay constant f_π. When we turn on the $SU(2)_L \otimes U(1)_Y$ electroweak interaction, the electroweak gauge bosons couple to the axial currents and acquire masses of order $\sim g f_\pi$. The massless pions thus disappear from the physical spectrum, having become the longitudinal components of the weak gauge bosons. Unfortunately, the mass acquired by the intermediate bosons is far smaller than required for a successful low-energy phenomenology; it is only[21] $M_W \sim 30\ \mathrm{MeV}/c^2$.

The minimal technicolor model of Weinberg[22] and Susskind[23] transcribes the same ideas from QCD to a new setting. The technicolor gauge group is taken to be $SU(N)_{TC}$ (usually $SU(4)_{TC}$), so the gauge interactions of the theory are generated by

$$SU(4)_{TC} \otimes SU(3)_c \otimes SU(2)_L \otimes U(1)_Y \ . \tag{9.2}$$

The technifermions are a chiral doublet of massless color singlets

$$\begin{pmatrix} U \\ D \end{pmatrix}_L \quad U_R,\ D_R \ . \tag{9.3}$$

With the electric charge assignments $Q(U) = \frac{1}{2}$ and $Q(D) = -\frac{1}{2}$, the theory is free of electroweak anomalies. The ordinary fermions are all technicolor singlets.

In analogy with our discussion of chiral symmetry breaking in QCD, we assume that the chiral TC symmetry is broken,

$$SU(2)_L \otimes SU(2)_R \otimes U(1)_V \to SU(2)_V \otimes U(1)_V \ . \tag{9.4}$$

Three would-be Goldstone bosons emerge. These are the technipions

$$\pi_T^+, \quad \pi_T^0, \quad \pi_T^-, \tag{9.5}$$

for which we are free to *choose* the technipion decay constant as

$$F_\pi = \left(G_F\sqrt{2}\right)^{-1/2} = 247\ \mathrm{GeV}\ . \tag{9.6}$$

This amounts to choosing the scale on which technicolor becomes strong. When the electroweak interactions are turned on, the technipions become the longitudinal components of the intermediate bosons, which acquire masses

$$\begin{aligned} M_W^2 &= g^2F_\pi^2/4 = \frac{\pi\alpha}{G_F\sqrt{2}\sin^2\theta_W} \\ M_Z^2 &= (g^2+g'^2)\,F_\pi^2/4 = M_W^2/\cos^2\theta_W \end{aligned} \tag{9.7}$$

that have the canonical Standard Model values, thanks to our choice (9.6) of the technipion decay constant.

Technicolor shows how the generation of intermediate boson masses could arise without fundamental scalars or unnatural adjustments of parameters. It thus provides an elegant solution to the naturalness problem of the Standard Model. However, it has a major deficiency: it offers no explanation for the origin of quark and lepton masses, because no Yukawa couplings are generated between Higgs fields and quarks or leptons.

A possible approach to the problem of quark and lepton masses is suggested by "extended technicolor" models and their modern extensions, "walking technicolor" models.[24] Technicolor implies a number of spinless technipions with masses below the technicolor scale of about 1 TeV. Some of these, the color singlet, technicolor singlet particles, should be relatively light. The colored technipions and technivector mesons may just be accessible to experiments at the Tevatron, but a thorough investigation awaits experiments on the 1-TeV scale.

10 Concluding remarks

We have recognized the significance of the 1-TeV scale—the realm of electroweak symmetry breaking—for nearly two decades. Through the development of superconducting magnets, and thanks to the experience gained in operating high-energy $\bar{p}p$ colliders at CERN and Fermilab and the evolution of detector architecture from Mark I at SPEAR up through UA1 and UA2 at CERN and CDF and D0 at Fermilab, we now have in hand the technical means to begin our assault on this frontier of our understanding. In addition, we look forward to a rich program of search and to testing our current understanding. The "bread-and-butter" physics of the LHC, including the detailed study of top quarks and B-particles, possibly extending to the study of CP violation, is an exciting prospect in its own right.

Avanti!

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