Higgs Physics: An Historical Perspective

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Abstract

“Weakly-coupled” and “strongly-coupled” models of electroweak symmetry breaking are introduced by analogy with the Fermi theory of the weak interaction and the low-energy interaction of pions, respectively. The implications of these two classes of models for colliders beyond the LHC and NLC are discussed.

*Presented at the Symposium on Future High Energy Colliders, Institute for Theoretical Physics, University of California at Santa Barbara, October 21-25, 1996.
As we have heard repeatedly at this meeting, uncovering the mechanism which breaks the electroweak symmetry will be a crucial step forward in our quest to understand nature at a deeper level. Electroweak symmetry breaking is the target of future accelerators, and we will hear a great deal about the phenomenology of this physics in talks on the LHC, NLC and \( \mu^+\mu^- \) colliders. I have therefore chosen topics with an eye towards minimizing overlap with other talks.

As my title indicates, I discuss electroweak symmetry breaking from an historical perspective. Since we have not yet discovered the physics of electroweak symmetry breaking, this requires some imagination. I begin with a very brief history of the weak interaction. I then discuss the two categories of electroweak symmetry breaking, usually referred to as “weak coupling” and “strong coupling.” I next discuss the phenomenological implications of these two categories of electroweak symmetry breaking for accelerators beyond the LHC and NLC. I conclude with a few historical remarks.

1 A brief “history” of the weak interaction

This section is intended to recall some of the milestones in the development of the standard model of the electroweak interaction. It is too sketchy to be a proper history. Furthermore, it is more a history of the way things could have gone, rather than the way they actually went. For a proper historical account, see Ref. [1].

The study of the weak interaction began 100 years ago, with the discovery of beta decay by H. Becquerel. It was not recognized at the time that this was due to a new force, however. The first “modern” theory of the weak interaction was due to Fermi (1933) [2], and involved a four-fermion interaction Lagrangian

\[
L = G_F \bar{\psi} \gamma^\mu \psi \bar{\psi} \gamma^\mu \psi
\]

where \( G_F \) is the Fermi constant and the fermion fields are those of the proton, neutron, electron, and electron neutrino.\(^b\) The discovery of the muon in 1947 allowed the study of the weak interaction in another context, and it was soon realized that the weak interaction is the same for electrons and muons. The universality of the weak interaction suggested that it is mediated by a gauge boson, in analogy with quantum electrodynamics, and that the electron and muon have the same weak “charge.” However, gauge bosons are exactly massless, while the hypothetical weak gauge boson is necessarily massive, since it yields the Fermi theory at energies much less than the weak-boson mass. This obstacle was surmounted in 1967 by Weinberg and Salam, who argued that the weak interaction is indeed a gauge theory, but with the gauge symmetry spontaneously broken, such that the weak gauge boson acquires a mass \(^3\).

We now jump to 1996, where we have a beautiful theory of the electroweak (and strong) interactions acting on three generations of quarks and leptons. It is fair to say that the gauge interactions are understood, and that the masses of the weak gauge bosons, the \( W \) and \( Z \), are understood to be a consequence of electroweak symmetry breaking. However,
the mechanism responsible for electroweak symmetry breaking is unknown. This mechanism is also responsible, at least in part, for the fermion masses and the Cabibbo-Kobayashi-Maskawa matrix (including CP violation), so its complete elucidation is essential to our quest to understand nature at a deeper level.

Although we do not know what the electroweak-symmetry-breaking mechanism is, we know that it involves new particles and new forces. It is possible that it will take another 100 years before we completely understand the nature of these particles and forces.

2 Electroweak Symmetry Breaking

2.1 Weak coupling

To explain what is meant by weakly-coupled electroweak symmetry breaking, I draw an analogy with the Fermi theory, which is the low-energy limit of a weakly-coupled theory, namely the electroweak gauge theory.

2.1.1 Fermi theory

Consider the calculation of a four-fermion amplitude in the Fermi theory of the weak interaction, as shown in Fig. 1. We use the interaction Lagrangian, Eq. (1), to calculate the amplitude perturbatively. We can use dimensional analysis to derive the dependence of the amplitude on the typical energy in the process, $E$. The Fermi constant, $G_F$, has units of inverse energy squared. Since the amplitude is dimensionless, the leading term in the amplitude, from the diagram in Fig. 1(a), is proportional to $G_F E^2$.

At next order in perturbation theory, one has the one-loop diagram in Fig. 1(b). Using dimensional analysis, we see that this diagram is proportional to $G_F^2 E^4$. Thus we are performing a perturbative expansion in powers of the dimensionless quantity $G_F E^2$.

This statement is in fact correct, but the argument is somewhat more subtle, because the loop integration is ultraviolet divergent. The modern attitude towards this divergence...
is encompassed by the idea of an “effective field theory” [4, 5]. Let’s say we only believe the Fermi theory up to some energy $\Lambda \gg E$. We restrict the loop integration to momenta less than $\Lambda$, and profess ignorance about what happens for momenta greater than $\Lambda$. However, we cannot simply throw away the contribution from momenta greater than $\Lambda$. Instead, we parameterize it by adding additional terms to the interaction Lagrangian of the form

$$\mathcal{L} = c G_F^2 \partial^\nu (\bar{\psi} \gamma^\mu \psi) \partial_\nu (\bar{\psi} \gamma_\mu \psi) + \ldots$$

(2)

which are characterized by having two derivatives, and an unknown dimensionless coefficient $c$. These terms yield a tree-level contribution to the amplitude proportional to $c G_F^2 E^4$, indicated in Fig. 1(c). Thus the next-to-leading-order amplitude is given schematically by

$$A = G_F E^2 + G_F^2 E^4 \ln \Lambda + c G_F^2 E^4$$

(3)

where the ultraviolet divergence of the loop diagram is evidenced by the dependence of the amplitude on $\ln \Lambda$. Combining the last two terms using $c' \equiv c + \ln \Lambda$ yields

$$A = G_F E^2 + c' G_F^2 E^4$$

(4)

where $c'$ is to be taken from experiment. Thus the amplitude is indeed an expansion in powers of the dimensionless quantity $G_F E^2$, despite the ultraviolet divergence.

At low energies, $E \ll G_F^{-1/2}$, only the first few terms in the expansion are numerically important, and the amplitude depends on just a few coefficients which must be extracted from experiment. However, for $E \sim G_F^{-1/2}$, every term in the expansion is equally important, and the amplitude depends on an infinite number of unknown coefficients. The theory therefore loses all predictive power. It is sometimes said that the theory “breaks down” or becomes “strongly-interacting,” but in fact the theory simply becomes useless.

There is a quick way to estimate the critical energy at which the theory becomes unpredictable, which also allows us to get the numerical factors straight. Unitarity implies that the partial waves of a two-particle scattering amplitude have a real part which does not exceed 1/2 in magnitude. If the leading term in the expansion were to greatly exceed this bound, the higher-order terms in the expansion would have to be just as large as the leading term in order to restore unitarity. But if that were the case, then every term in the expansion is equally important, and the theory is unpredictable. So we can estimate the critical energy by imposing the unitarity bound on the leading-order amplitude. This yields

$$E_{\text{critical}} \approx \left( \frac{\sqrt{2} \pi}{G_F} \right)^{1/2} \approx 600 \text{ GeV}.$$ 

(5)

Thus the Fermi theory loses predictive power for $E \sim 600 \text{ GeV}$.

It is a logical possibility that physics simply becomes very complicated at energies in excess of the critical energy. However, history has taught us to expect the opposite: the theory becomes simpler at higher energy. In the case of the Fermi theory, the four-fermion

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\footnote{This is reminiscent of a quote from the early days of renormalization theory: “Just because something is infinite does not mean it is zero!” [3]}
Figure 2: The Fermi theory of the weak interaction is the low-energy approximation to a gauge theory.

interaction is replaced by the exchange of a weak gauge boson, as shown in Fig. 2. The amplitude is proportional to

$$A \sim g^2 \frac{E^2}{E^2 - M_W^2}$$

(6)

where $g$ is the gauge coupling, and the denominator is from the $W$ propagator. Since

$$G_F \sim \frac{g^2}{M_W^2}$$

(7)

the amplitude reduces to that of the Fermi theory for $E \ll M_W$. However, for $E \gg M_W$, the amplitude is proportional to $g^2$. Hence the expansion parameter is a (small) constant, and the theory is predictive for all energies.

In the case of the Fermi theory, new physics, in the form of the $W$ boson, enters before the critical energy. Formally,

$$M_W^2 \sim \frac{g^2}{4\pi} E_{critical}^2$$

(8)

using Eqs. (5) and (7). The fact that the new physics enters before the critical energy can thus be related to the fact that $g^2/4\pi < 1$, i.e., that the theory is weakly coupled.

2.1.2 Electroweak theory

In the previous section, the $W$ boson was responsible for regulating the growth of the four-fermion amplitude at high energy. Now consider the scattering of the $W$ bosons themselves, in particular longitudinal (helicity-zero) $W$ bosons, as shown in Fig. 3. The top row of Feynman diagrams depicts the gauge interactions responsible for the scattering. The resulting amplitude is proportional to $G_F E^2$, just as in the case of the Fermi theory. Using unitarity, we estimate the critical energy to be

$$E_{critical} \approx \left( \frac{4\sqrt{2}\pi}{G_F} \right)^{1/2} \approx 1.2 \text{ TeV}.$$  

(9)

Thus the electroweak theory loses predictive power for $E \sim 1.2 \text{ TeV}$.

\[ \text{More precisely, using } G_F/\sqrt{2} = g^2/8M_W^2. \]
By analogy with the Fermi theory, we might expect new physics, in the form of a new particle, to regulate the growth of the amplitude. This is exactly what happens in the standard Higgs model. The Higgs boson gives rise to additional Feynman diagrams, shown in the second row in Fig. 3. These diagrams cancel the terms proportional to $G_F E^2$, leaving behind an amplitude proportional to $G_F m_H^2 \sim \lambda$, where $\lambda$ is the Higgs self-coupling. Thus the expansion parameter is a (small) constant, and the theory is predictive for all energies. Higgs-higgs scattering is also proportional to $\lambda$, so the theory is complete.

The formal relation between the Higgs mass and the critical energy is

$$m_H^2 \sim \frac{\lambda}{4\pi} E_{\text{critical}}^2 < E_{\text{critical}}^2$$  \hspace{1cm} (10)$$

where the last relation relies on $\lambda/4\pi < 1$, known from non-perturbative studies of the standard Higgs model \[9\]. Thus the new physics, namely the Higgs boson, enters before the critical energy because the theory is weakly coupled. In fact, it has been shown that the standard Higgs model, and variations of it (such as two Higgs doublets, as used in supersymmetric models), are the only weakly-coupled models of electroweak symmetry breaking \[10\].

### 2.2 Strong coupling

To explain what is meant by strongly-coupled electroweak symmetry breaking, I draw an analogy with low-energy pion physics, which is the low-energy limit of a strongly-coupled theory, namely QCD.

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\[8\]Ignoring the running of $\lambda$, to be discussed in a later footnote.
2.2.1 Pion theory

The low-energy interaction of pions is dictated by the fact that they are the (approximate) Goldstone bosons of spontaneously-broken chiral symmetry. This is embodied by an effective field theory of pions called “chiral perturbation theory” [4]. The leading interaction of pions is

\[ \mathcal{L} = G_\pi \pi \cdot \partial^\mu \pi \pi \cdot \partial_\mu \pi \]  (11)

where \( G_\pi = 1/(2f_\pi^2) \), with \( f_\pi \) the pion decay constant. The coupling \( G_\pi \) has dimensions of inverse energy squared, in analogy with the Fermi constant, \( G_F \).

The leading-order amplitude for \( \pi\pi \) scattering is proportional to \( G_\pi E^2 \). Using unitarity to estimate the critical energy at which the theory loses predictive power yields

\[ E_{\text{critical}} \sim \left( \frac{4\pi}{G_\pi} \right)^{1/2} \approx 450 \text{ MeV} . \] (12)

What actually happens in nature in \( \pi\pi \) scattering at 450 MeV? One begins to encounter a plethora of new particles, beginning with the \( \sigma \) meson. The theory of pion interactions becomes complicated above 450 MeV, and loses all predictive power. Unlike the case of a weakly-coupled theory, there is no new particle (the analogue of a \( W \) or a Higgs boson) that restores predictivity to the theory.

Although the pion theory becomes complicated above 450 MeV, we have learned that nature is nevertheless simple: there is a new description of physics in terms of quarks and gluons, interacting via QCD. At low energies this theory is strongly-interacting, and gives rise to all the complications of hadron physics. But even at low energies the theory itself is simple, in the sense that it is described by a Lagrangian with just a few terms in it, and the only parameters are the strong coupling constant and the quark masses.

2.2.2 Electroweak theory

Now consider \( WW \) scattering near the critical energy of 1.2 TeV, and imagine that nature does not provide a Higgs boson to regulate the growth of the amplitude with energy. What will happen? There is no answer to this question within the electroweak theory; it simply becomes unpredictive. However, based on our experience with pion physics, we might expect that there is a deeper, simpler theory, which is hidden from view by virtue of the fact that it is strongly interacting, and therefore manifests itself in a complicated way. This would mean that there are new particles in nature, experiencing a new strong interaction.

For example, one can imagine that this simpler theory is analogous to QCD. This is the idea behind the so-called Technicolor theory [12, 13]. In that case, one would encounter the analogues of the \( \sigma \), \( \rho(770) \), etc., at energies above 1.2 TeV. These “Technimesons” would be made from strongly-interacting “Techniquarks,” bound together by “Technigluons.” This class of models is reviewed by Appelquist at this symposium.

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\(^{1}\)After a twenty-year absence, the \( \sigma \) meson has once again been recognized by the Particle Data Group, but they are hedging on its mass: \( f_0(400 - 1400) \) [1].

\(^{2}\)One might be tempted to interpret the \( \sigma \) as the Higgs boson of the pion theory. This interpretation is invalidated by the \( \rho(770) \), which is comparable in mass to the \( \sigma \). A true Higgs theory would have no such particle.
In this section I make some observations about the phenomenological implications of electroweak symmetry breaking for colliders with energy in excess of the LHC ($\sqrt{s} = 14$ TeV) and the NLC ($\sqrt{s} = 0.5 - 1.5$ TeV). I begin where I left off, with the case of strongly-coupled electroweak symmetry breaking.

3 Phenomenology beyond the LHC/NLC

Let’s continue our analogy with pion physics. To study the strong interaction at low energy, one performs scattering experiments involving pions. For example, one can perform $\pi\pi$ scattering experiments, or study the coupling of the pion to virtual photons, as depicted in Fig. 4(a). The analogues of these processes for the electroweak interaction are shown in Fig. 4(b), where the longitudinal $W$ bosons replace the pions. The incident fermions could be either quarks and antiquarks (LHC) or electrons and positrons (NLC). These colliders will be capable of probing these processes at energies of about 1 TeV, comparable to the critical energy of 1.2 TeV at which the theory loses predictive power.

Recall that the critical energy for $\pi\pi$ scattering is about 450 MeV. How much would we know about the strong interaction if we had data from $\pi\pi$ scattering and $\gamma^* \to \pi\pi$ at energies only up to 450 MeV? The answer is very little. We might know about the $\sigma$, but we would see only the low-energy tail of the $\rho$, and the heavier mesons would be completely out of sight. More importantly, we would not know that the mesons are composed of strongly-interacting...
quarks, interacting via QCD.

The moral is that if the mechanism of electroweak symmetry breaking is indeed strongly-coupled, it will likely require energies greatly in excess of 1 TeV for the complete elucidation of this physics. This implies the need for colliders beyond the LHC and NLC. Although these machines are likely to tell us something about strongly-coupled electroweak symmetry breaking, it is hard to imagine they will be able to tell us everything about it.

3.2 Weak coupling

The case of weakly-coupled electroweak symmetry breaking is apparently in stark contrast with that of the strongly-coupled case. For example, the standard Higgs model has a Higgs boson with mass less than 700 GeV [9], and nothing else. The LHC and NLC are both capable of discovering this particle. Once discovered, and its coupling to itself and other particles measured, we have learned everything there is to learn about electroweak symmetry breaking. Can it really be this simple?

Probably not. Although the standard Higgs model is predictive at all energies, it suffers from another disease - it is “unnatural” [14]. We know that the Higgs field acquires a vacuum-expectation value of \( v = (\sqrt{2}G_F)^{-1/2} \approx 250 \text{ GeV} \). The diagram in Fig. 5(a) is one of the one-loop corrections to the vacuum-expectation value, from a Higgs loop. There are similar one-loop corrections from loops of weak bosons and fermions. All these one-loop diagrams share the feature that they are quadratically divergent [13]. Let us regard the standard Higgs model as being valid up to some energy \( \Lambda \). The relation between the bare vacuum-expectation value, \( v_0 \), and the actual vacuum-expectation value, \( v \), may be approximated by cutting off the loop integration at momenta of order \( \Lambda \):

\[
v^2 = v_0^2 + O \left( \frac{\Lambda^2}{(4\pi)^2} \right)
\]

where the factor \((4\pi)^2\) is the usual factor which arises from loop diagrams. There is also an (unknown) contribution from momenta greater than \( \Lambda \), but we will assume that it does not conspire to cancel the contribution from momenta less than \( \Lambda \). If \( \Lambda >> v \), then \( v_0 \) must be tuned to almost exactly cancel the one-loop contribution to the vacuum-expectation value. Such a fine tuning is unnatural. Instead, it is natural to expect \( \Lambda/(4\pi) \leq v \). Thus if nature makes use of the standard Higgs model, we anticipate that it is replaced by a more fundamental theory at energy \( \Lambda \leq 4\pi v \approx 3 \text{ TeV} \), regardless of the Higgs mass.

\[\text{It has been argued that just such a cancellation occurs if the underlying theory, to which the standard Higgs model is a low-energy approximation, is conformally invariant } [15]. \text{ However, a realistic example of such a theory has yet to be constructed.}\]

\[\text{A separate argument can be given for the incompleteness of the standard Higgs model. The Higgs self-coupling, } \lambda, \text{ is a running coupling, and increases with energy (for sufficiently-large Higgs mass), eventually blowing up. New physics must intervene before this occurs. This can be used to place an upper bound on the Higgs mass for a given energy scale of the new physics } [16]. \text{ For example, if the new physics is at the Planck scale, there is an upper bound on the Higgs mass of about } 200 \text{ GeV } [17]. \text{ However, naturalness implies that there is new physics by at least } 3 \text{ TeV, so it is not realistic to imagine that there is no new physics until the Planck scale. Imposing the condition that the Higgs coupling not blow up below } 3 \text{ TeV yields an upper bound on the Higgs mass of only about } 700 \text{ GeV.}\]
Thus we see that weakly-coupled electroweak symmetry breaking may not be so different from strongly-coupled electroweak symmetry breaking. The strongly-coupled approach involves new physics at the TeV scale, while the weakly-coupled approach suggests new physics (beyond the standard Higgs model) by at least the TeV scale. It could be that this new physics lies well below the TeV scale, in which case it may be possible to discover all of it with the LHC/NLC. But if it really lies at the TeV scale, it will likely require colliders beyond the LHC/NLC for its complete elucidation.

A well-known example of new physics in the weakly-coupled scenario is supersymmetry. If nature is supersymmetric, every particle is accompanied by a superpartner. The one-loop correction to the Higgs vacuum-expectation value in Fig. 5(a) is accompanied by a second diagram, shown in Fig. 5(b), in which the Higgs loop is replaced by a Higgsino loop. The quadratic divergences of each diagram cancel, and the theory becomes natural. The same sort of cancellation occurs for loops of weak bosons and fermions and their superpartners.

If nature is supersymmetric, then supersymmetry must be broken, since we don’t observe the superpartners of the known particles. Let’s refer to the mass scale of the superpartners as $M_{SUSY}$. It $M_{SUSY}$ were much larger than $v$, the quadratic divergence would reappear, cut off only when the momenta reach $M_{SUSY}$:

$$v^2 = v_0^2 + O \left( \frac{M_{SUSY}^2}{(4\pi)^2} \right).$$

Thus the model is only natural if $M_{SUSY} < 4\pi v \approx 3 \text{ TeV}$. More careful estimates yield a somewhat lower scale [18].

If $M_{SUSY}$ is at the TeV scale, it may require colliders beyond the LHC/NLC to discover all the superpartners and measure their couplings. However, if $M_{SUSY}$ is a few hundred GeV, then all the physics of supersymmetry may be within reach of the LHC/NLC. This is the only natural scenario I know of in which the entire physics of electroweak symmetry breaking is elucidated by these colliders. However, even in this scenario, there may be motivation for higher-energy colliders. The physics of supersymmetry breaking may lie at the 10 TeV scale, as in the class of models in which dynamical supersymmetry breaking is communicated to the standard model via new gauge interactions [19]. It is possible that the elucidation of the physics of electroweak symmetry breaking will reveal to us the scale of supersymmetry breaking.
It is also possible that the new physics which subsumes the standard Higgs model at the TeV scale is strongly-interacting, and has nothing to do with supersymmetry \[20\]. This class of models is similar to the strongly-coupled electroweak-symmetry-breaking scenario, with the exception that there is a Higgs boson. The complete elucidation of these models would likely require colliders beyond the LHC/NLC, as in the strongly-coupled models.

4 Concluding historical remarks

The history of the weak interaction was punctuated by periods of confusion, followed by clarification, which ultimately led to a beautiful theory. For example, the early measurements of beta decay indicated that the emitted electron was monoenergetic. It took twenty years to establish that the electron is emitted with a continuous spectrum of energies. This was the first evidence for the existence of the neutrino, which was a vital ingredient in Fermi’s theory of the weak interaction \[21\].

The same sequence of events is likely to occur for the electroweak-symmetry-breaking mechanism. In fact, we are already in the first stage, the period of confusion. There are already several experimental results which have been interpreted as a Higgs boson \[22, 23, 24, 25, 26, 27, 28\]. These interpretations have been made either in the context of a multi-Higgs model, or with additional new physics – none has had an interpretation in terms of the standard Higgs model with no new physics.

A particularly noteworthy example is Ref. \[26\], which was an attempt to interpret the $\zeta(8.3)$ as a Higgs boson with enhanced coupling to $b$ quarks.\footnote{The $\zeta(8.3)$ was conjectured to be responsible for a monoenergetic photon signal in $\Upsilon$ decay, via $\Upsilon \rightarrow \zeta \gamma$. The experiment turned out to be erroneous.} This would imply an enhanced coupling to muons, which led the authors to conclude that if their interpretation were correct, “the case for construction of a dedicated muon collider for Higgs boson studies will become as compelling as it is technically feasible.” The idea of using a muon collider for resonant Higgs production (even the standard Higgs) has lately resurfaced and been considerably refined \[29\].

My final historical remark concerns our ability to foresee the physics of the future. There are many examples of incorrect theories and prognostications, many more than correct ones, and these are often held up as examples of our inability to predict what will be found in the next generation of colliders. While I have some sympathy for this point, I also feel it can be overstated. A well-known example is the top-quark mass; predictions ranged from 15 GeV (just above the reach of PEP) to 230 GeV, and everything in between. But perhaps this makes us lose sight of the real achievement: we knew the top quark existed long before it was discovered, based entirely on indirect evidence.

Today many feel that there is indirect evidence for weak-scale supersymmetry. Although the evidence is not as compelling as it was for the top quark, an advocate might make the following statement:

“There can be no two opinions about the practical utility of supersymmetry, but until clear experimental evidence for the existence of supersymmetry can be
obtained, supersymmetry must remain purely hypothetical. Failure to detect any
evidence of supersymmetry is no evidence against its existence.”

In fact, this quote is paraphrased from a 1937 meeting of the Royal Society, except I have
substituted “supersymmetry” for “the neutrino” throughout [30].

Acknowledgements

I am grateful for conversations with D. Dicus, R. Leigh, and W. Marciano, and for
assistance from Z. Sullivan. This work was supported in part by the Department of Energy
under Grant No. DE-FG02-91ER40677 and by the National Science Foundation under Grant
No. PHY94-07194.

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