

The State of the Standard Model

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Dedicated to the Memory of Sam Treiman

Abstract. I quickly review the successes of quantum chromodynamics. Then I assess the current state of the electroweak theory, making brief comments about the search for the Higgs boson and some of the open issues for the theory. I sketch the problems of mass and mass scales, and point to a speculative link between the question of identity and large extra dimensions. To conclude, I return to QCD and the possibility that its phase structure might inform our understanding of electroweak symmetry breaking.

APPRECIATION

Sam Treiman left us a few days ago. With his deep respect for great principles, his acute understanding of field theory, and his joyful engagement with experiment, Sam was a model and a guide to many of us for whom the goal—the joy—of theoretical physics is not to gather up a hoard of shiny theorems, but to learn to read Nature’s secrets. He was mentor to many of our most valued theoretical colleagues, and the course on particle physics he described as “shamelessly phenomenological” brought culture to generations of Princeton students. He was a natural man—unpretentious and a little ruffled—who took wry pleasure in seeing the self-important deflated. A graceful writer himself [1,2], he was a writer’s ideal reader: curious, appreciative, eager to share an original idea or a provocative passage from his reading, and ever ready to recommend a good book.² Sam confessed to me, without remorse, that he coined “The Standard Model,” a curiously flat name for the marvelous theory of quarks and leptons that he helped to build. I forgave him then, and I forgive him now—but we still need a better name!

¹) Fermilab is operated by Universities Research Association Inc. under Contract No. DE-AC02-76CH03000 with the United States Department of Energy.

²) His last recommendation to me was the luminous translations of *The Iliad* and *The Odyssey* by Robert Fagles [3].

OUR PICTURE OF MATTER

Some twenty-five years after the November Revolution, we base our understanding of physical phenomena on the identification of a few constituents that seem elementary at the current limits of resolution of about 10^{-18} m, and a few fundamental forces. The constituents are the pointlike quarks

$$\begin{pmatrix} u \\ d \end{pmatrix}_L \quad \begin{pmatrix} c \\ s \end{pmatrix}_L \quad \begin{pmatrix} t \\ b \end{pmatrix}_L \quad (1)$$

and leptons

$$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L \quad \begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L \quad \begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}_L \quad (2)$$

with strong, weak, and electromagnetic interactions specified by $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ gauge symmetries.

This concise statement of the standard model invites us to consider the agenda of particle physics today under four themes. *Elementarity.* Are the quarks and leptons structureless, or will we find that they are composite particles with internal structures that help us understand the properties of the individual quarks and leptons? *Symmetry.* One of the most powerful lessons of the modern synthesis of particle physics is that symmetries prescribe interactions. Our investigation of symmetry must address the question of which gauge symmetries exist (and, eventually, why). We must also understand how the electroweak symmetry³ is hidden. The most urgent problem in particle physics is to complete our understanding of electroweak symmetry breaking by exploring the 1-TeV scale. *Unity.* We have the fascinating possibility of gauge coupling unification, the idea that all the interactions we encounter have a common origin and thus a common strength at suitably high energy. Next comes the imperative of anomaly freedom in the electroweak theory, which urges us to treat quarks and leptons together, not as completely independent species. Both ideas are embodied in unified theories of the strong, weak, and electromagnetic interactions, which imply the existence of still other forces—to complete the grander gauge group of the unified theory—including interactions that change quarks into leptons. The self-interacting quanta of non-Abelian theories and supersymmetry both hint that the traditional distinction between force particles and constituents might give way to a unified understanding of all the particles. *Identity.* We do not understand the physics that sets quark masses and mixings. Although experiments are testing the idea that the phase in the quark-mixing matrix lies behind the observed CP violation, we do not know what determines that phase. The accumulating evidence for neutrino oscillations presents us with a new embodiment of these puzzles in the lepton sector. At bottom, the question of identity is very simple to state: What makes an electron an electron, a neutrino a neutrino, and a top quark a top quark?

³⁾ and, no doubt, others—including the symmetry that brings together the strong, weak, and electromagnetic interactions.

QCD IS PART OF THE STANDARD MODEL

The quark model of hadron structure and the parton model of hard-scattering processes have such pervasive influence on the way we conceptualize particle physics that quantum chromodynamics, the theory of strong interactions that underlies both, often fades into the background when the standard model is discussed. I want to begin my state of the standard model report with the clear statement that QCD is indeed part of the standard model, and with the belief that understanding QCD may be indispensable for deepening our understanding of the electroweak theory.

Quantum chromodynamics is a remarkably simple, successful, and rich theory of the strong interactions.⁴ The perturbative regime of QCD exists, thanks to the crucial property of asymptotic freedom, and describes many phenomena in quantitative detail. The strong-coupling regime controls hadron structure and gives us our best information about quark masses.

The classic test of perturbative QCD is the prediction of subtle violations of Bjorken scaling in deeply inelastic lepton scattering. As an illustration of the current state of the comparison between theory and experiment, I show in Figure 1 the singlet structure function $F_2(x, Q^2)$ measured in νN charged-current interactions by the CCFR Collaboration at Fermilab. The solid lines for $Q^2 \gtrsim (5 \text{ GeV}/c)^2$ represent QCD fits; the dashed lines are extrapolations to smaller values of Q^2 . As we see in this example, modern data are so precise that one can search for small departures from the QCD expectation.

Perturbative QCD also makes spectacularly successful predictions for hadronic processes. I show in Figure 2 that pQCD, evaluated at next-to-leading order using the program JETRAD, accounts for the transverse-energy spectrum of central jets produced in the reaction

$$\bar{p}p \rightarrow \text{jet}_1 + \text{jet}_2 + \text{anything} \quad (3)$$

over at least six orders of magnitude, at $\sqrt{s} = 1.8 \text{ TeV}$.⁵

The Q^2 -evolution of the strong coupling constant predicted by QCD, which in lowest order is

$$1/\alpha_s(Q^2) = 1/\alpha_s(\mu^2) + \frac{33 - 2n_f}{12\pi} \log(Q^2/\mu^2), \quad (4)$$

where n_f is the number of active quark flavors, has been observed within individual experiments [9,10] and by comparing determinations made in different experiments at different scales. A compilation of $1/\alpha_s$ determinations from many experiments, shown in Figure 3, exhibits the expected behavior.⁶ When evolved to a common

⁴) For a passionate elaboration of this statement, see Frank Wilczek's keynote address at PANIC '99, Ref. [4]. An authoritative portrait of QCD and its many applications appears in the monograph by Ellis, Stirling, and Webber, Ref. [5].

⁵) For a systematic review of high- E_T jet production, see Blazey and Flaughner, Ref. [8].

⁶) A useful plot of α_s vs. Q^2 appears as Figure 9.2 of the *Review of Particle Physics*, Ref. [11].

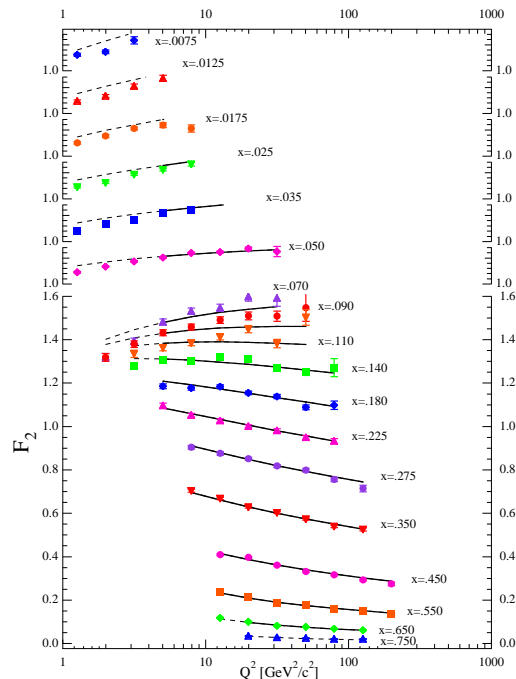


FIGURE 1. The structure function F_2 measured in νN interactions, from Ref. [6].

scale $\mu = M_Z$, the various determinations of α_s lead to consistent values, as shown in Figure 4.

Thanks to QCD, we have learned that the dominant contribution to the light-hadron masses is not the masses of the quarks of which they are constituted, but the energy stored up in confining the quarks in a tiny volume.⁷ Our most useful tool in the strong-coupling regime is lattice QCD. Calculating the light hadron spectrum from first principles has been one of the main objectives of the lattice program, and important strides have been made recently. In 1994, the GF11 Collaboration [13] carried out a quenched calculation of the spectrum (no dynamical fermions) that yielded masses that agree with experiment within 5–10%, with good understanding of the residual systematic uncertainties. The CP-PACS Collaboration centered in Tsukuba has embarked on an ambitious program that will soon lead to a full (unquenched) calculation. Their quenched results, along with those of the GF11 Collaboration, are presented in Figure 5 [14]. The gross features of the light-hadron spectrum are reproduced very well, but if you look with a critical eye (as the CP-PACS collaborators do), you will notice that the quenched light hadron spectrum systematically deviates from experiment. The K - K^* mass splitting is underestimated by about 10%, and the results differ depending on whether the strange-quark mass is fixed from the K mass or the ϕ mass. The forthcoming unquenched results should improve the situation further, and give us new insights into how well—and why!—the simple quark model works.

⁷⁾ An accessible essay on our understanding of hadron mass appears in Ref. [12].

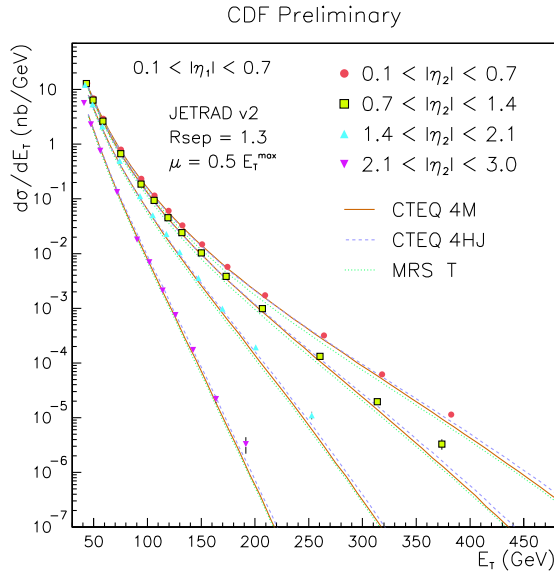


FIGURE 2. Cross sections measured at $\sqrt{s} = 1.8$ TeV by the CDF Collaboration for central jets (defined by $0.1 < |\eta_1| < 0.7$), with the second jet confined to specified intervals in the pseudorapidity η_2 [7]. The curves show next-to-leading-order QCD predictions based on the CTEQ4M (solid line), CTEQ4HJ (dashed line), and MRST (dotted line) parton distributions.

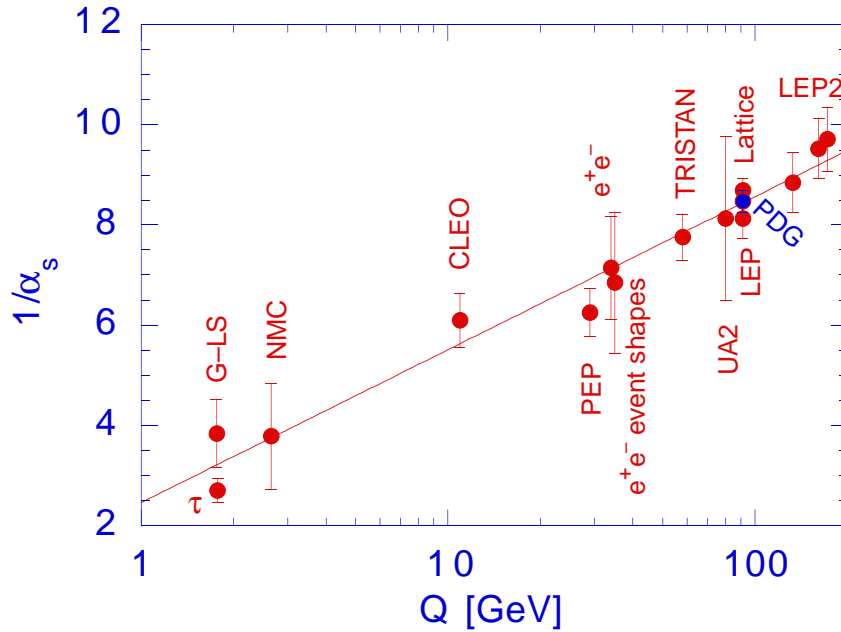


FIGURE 3. Determinations of $1/\alpha_s$, plotted at the scale μ at which the measurements were made. The line shows the expected evolution (4).

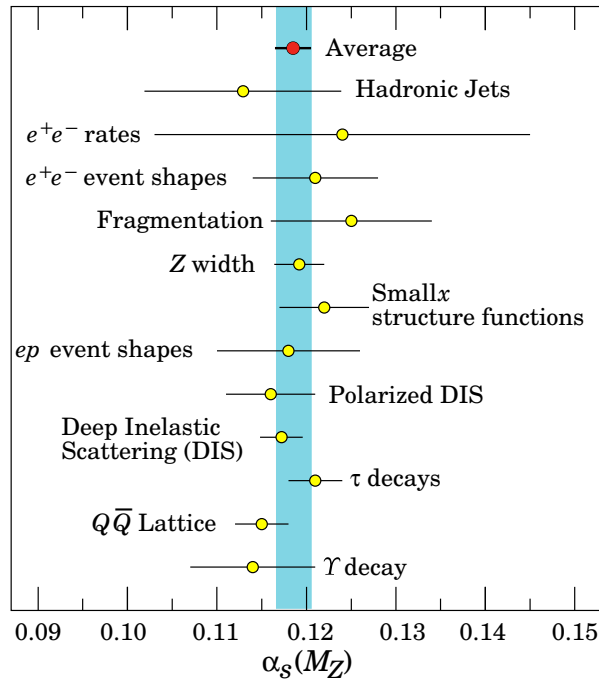


FIGURE 4. Determinations of $\alpha_s(M_Z)$ from several processes. In most cases, the value measured at a scale μ has been evolved to $\mu = M_Z$. Error bars include the theoretical uncertainties. From the *Review of Particle Physics* [11].

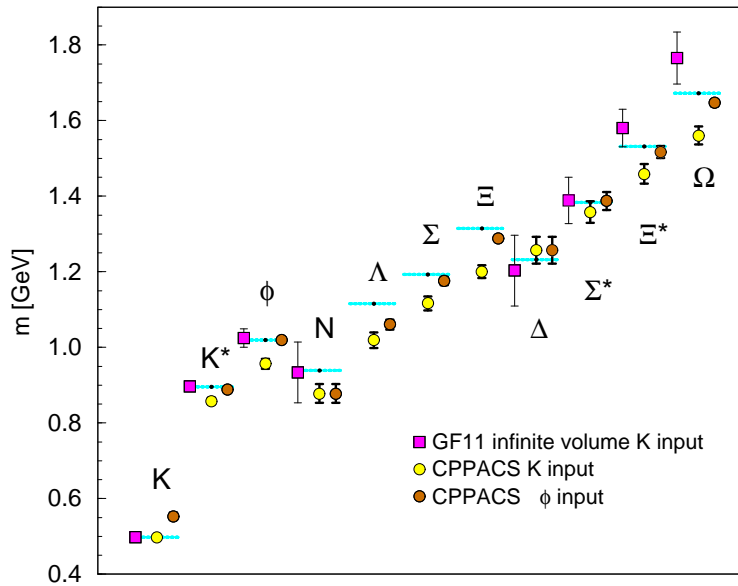


FIGURE 5. Final results of the CP-PACS Collaboration's quenched light hadron spectrum in the continuum limit [14]. Experimental values (horizontal lines) and earlier results from the GF11 Collaboration [13] are plotted for comparison.

THE $SU(2)_L \otimes U(1)_Y$ ELECTROWEAK THEORY

The electroweak theory is founded on the weak-isospin symmetry embodied in the doublets

$$\begin{pmatrix} u \\ d' \end{pmatrix}_L \quad \begin{pmatrix} c \\ s' \end{pmatrix}_L \quad \begin{pmatrix} t \\ b' \end{pmatrix}_L \quad \begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L \quad \begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L \quad \begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}_L \quad (5)$$

and weak-hypercharge phase symmetry, plus the idealization that neutrinos are massless.⁸ In its simplest form, with the electroweak gauge symmetry broken by the Higgs mechanism, the $SU(2)_L \otimes U(1)_Y$ theory has scored many qualitative successes: the prediction of neutral-current interactions, the necessity of charm, the prediction of the existence and properties of the weak bosons W^\pm and Z^0 . Over the past ten years, in great measure due to the beautiful experiments carried out at the Z factories at CERN and SLAC, precision measurements have tested the electroweak theory as a quantum field theory, at the one-per-mille level [16,17]. As an example of the insights precision measurements have brought us (one that mightily impressed the Royal Swedish Academy of Sciences), I show in Figure 6 the time evolution of the top-quark mass favored by simultaneous fits to many electroweak observables. Higher-order processes involving virtual top quarks are an important element in quantum corrections to the predictions the electroweak theory makes for many observables. A case in point is the total decay rate, or width, of the Z^0 boson: the comparison of experiment and theory shown in the inset to Figure 6 favors a top mass in the neighborhood of 180 GeV/ c^2 .

The comparison between the electroweak theory and a considerable universe of data is shown in Figure 7 [19], where the pull, or difference between the global fit and measured value in units of standard deviations, is shown for some twenty observables. The distribution of pulls for this fit, due to the LEP Electroweak Working Group, is not noticeably different from a normal distribution, and only a couple of observables differ from the fit by as much as about two standard deviations. This is the case for any of the recent fits, including two that I point to later as examples of what we may learn from global fits about physics beyond the standard model [20,21]. From fits of the kind represented here, we learn that the standard-model interpretation of the data favors a light Higgs boson. We will hear about these conclusions in more detail from Bill Marciano [22], and I will add some cautionary remarks shortly.

The beautiful agreement between the electroweak theory and a vast array of data from neutrino interactions, hadron collisions, and electron-positron annihilations at the Z^0 pole and beyond means that electroweak studies have become a modern arena in which we can look for new physics “in the sixth place of decimals.”⁹

⁸⁾ For a survey of the electroweak theory, with many references, see Ref. [15].

⁹⁾ In *The Odd Quantum*, Sam Treiman [2] quotes from the 1898–99 University of Chicago catalogue: “While it is never safe to affirm that the future of the Physical Sciences has no marvels in store even more astonishing than those of the past, it seems probable that most of the grand

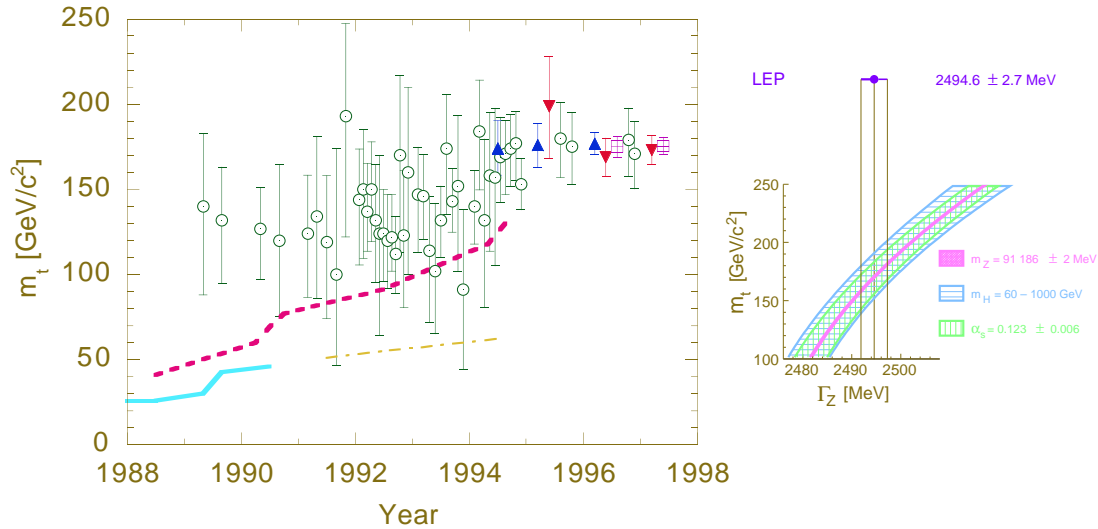


FIGURE 6. Indirect determinations of the top-quark mass from fits to electroweak observables (open circles) and 95% confidence-level lower bounds on the top-quark mass inferred from direct searches in e^+e^- annihilations (solid line) and in $\bar{p}p$ collisions, assuming that standard decay modes dominate (broken line). An indirect lower bound, derived from the W -boson width inferred from $\bar{p}p \rightarrow (W \text{ or } Z) + \text{anything}$, is shown as the dot-dashed line. Direct measurements of m_t by the CDF (triangles) and $D\bar{O}$ (inverted triangles) Collaborations are shown at the time of initial evidence, discovery claim, and at the conclusion of Run 1. The world average from direct observations is shown as the crossed box. For sources of data, see Ref. [11]. *Inset:* Electroweak theory predictions for the width of the Z^0 boson as a function of the top-quark mass, compared with the width measured in LEP experiments. (From Ref. [18]).

The Higgs Boson Search at the Tevatron

The most promising channel for Higgs-boson searches at the Tevatron will be the $b\bar{b}$ mode, for which the branching fraction exceeds about 50% throughout the region preferred by supersymmetry and the precision electroweak data. At the Tevatron, the direct production of a light Higgs boson in gluon-gluon fusion $gg \rightarrow H \rightarrow b\bar{b}$ is swamped by the ordinary QCD production of $b\bar{b}$ pairs. The high background in the $b\bar{b}$ channel means that special topologies must be employed to improve the ratio of signal to background and the significance of an observation. The high luminosities that can be contemplated for a future run argue that the associated-production reactions $\bar{p}p \rightarrow H(W, Z) + \text{anything}$ are plausible candidates for a Higgs discovery

underlying principles have been firmly established and that further advances are to be sought chiefly in the rigorous application of these principles to all the phenomena which come under our notice An eminent physicist has remarked that the future truths of Physical Science are to be looked for in the sixth place of decimals.” Future truths are to be found still in precision measurements, but the century we are leaving has repeatedly shown that Nature’s marvels are not limited by our imagination, and that exploration can yield surprises that completely change the way we think.

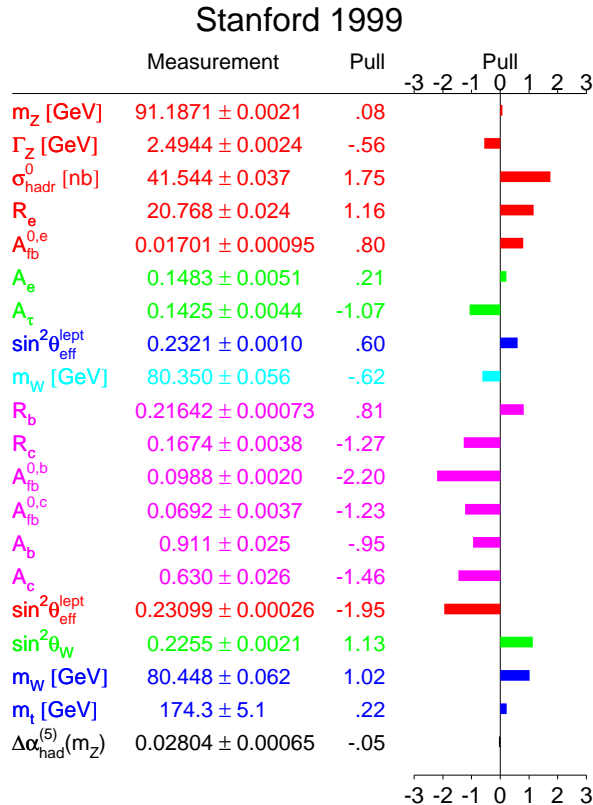


FIGURE 7. Precision electroweak measurements and the pulls they exert on a global fit to the standard model, from Ref. [19].

at the Tevatron.

The prospects for exploiting these reactions were explored in detail in connection with the Run 2 Supersymmetry / Higgs Workshop at Fermilab,¹⁰ and will be presented in detail at this meeting by John Hobbs [23]. Taking into account what is known, and what might conservatively be expected, about sensitivity, mass resolutions, and background rejection, these investigations show that it is unlikely that a standard-model Higgs boson could be observed in Tevatron Run 2. The prospects are much brighter for Run 2^{bis}. Indeed, the sensitivity to a light Higgs boson is what motivates the integrated luminosity of 30 fb^{-1} specified for Run 2^{bis}.

The detection strategy evolved in the Supersymmetry / Higgs Workshop involves combining the HZ and HW signatures and adding the data from the CDF and DØ detectors. Apart from the increase in luminosity, the key ingredient in the heightened sensitivity is a projected improvement in the $b\bar{b}$ mass resolution to 10%. Prospects are summarized in Figure 8, which shows as a function of the Higgs-boson mass the luminosity required for exclusion at 95% confidence level

¹⁰⁾ Work carried out in the context of the Tevatron Run 2 Supersymmetry/Higgs Workshop at Fermilab may be found at <http://fnth37.fnal.gov/susy.html>.

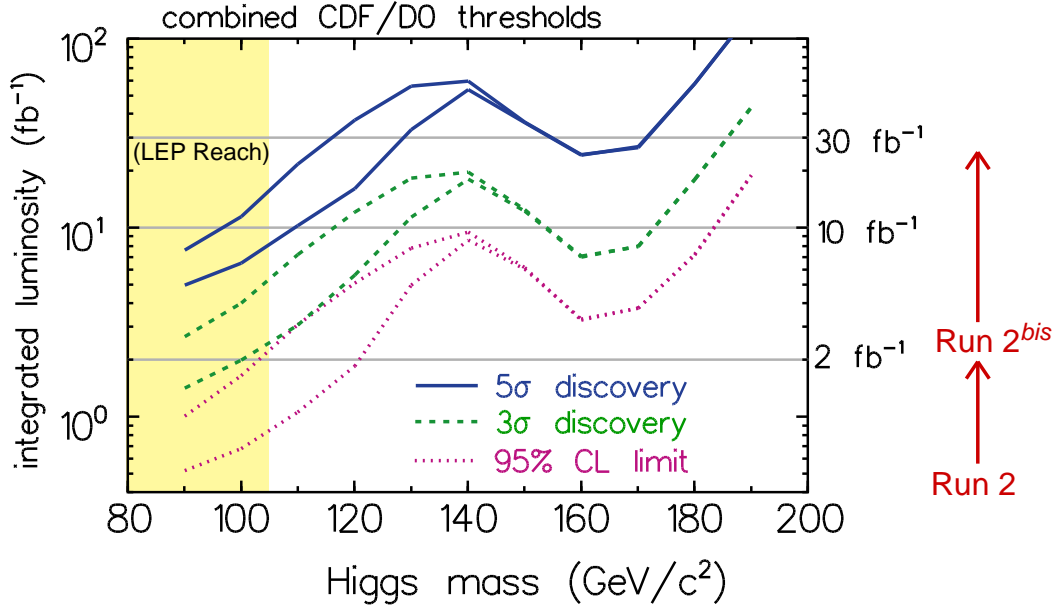


FIGURE 8. Integrated luminosity projected for the detection of a standard-model Higgs boson at the Tevatron Collider. In each pair of curves, the lower reach corresponds to a generic analysis, the upper reach to a neural-net optimization. See the talk by John Hobbs [23] for a detailed exposition of the assumptions.

(dotted lines), three-standard-deviation evidence (dashed lines), and five-standard-deviation discovery (solid lines). We see that an integrated luminosity of 2 fb^{-1} , expected in Run 2, is insufficient for a convincing observation of a standard-model Higgs boson with a mass too large to be observed at LEP 2. However, a 95% CL exclusion is possible up to about $125 \text{ GeV}/c^2$. On the other hand, about 10 fb^{-1} would permit detailed study of a standard-model Higgs boson discovered at LEP 2. If the Higgs boson lies beyond the reach of LEP 2, $M_H \gtrsim (105 - 110) \text{ GeV}/c^2$, then a $5\text{-}\sigma$ discovery will be possible in a future Run 2^{bis} of the Tevatron (30 fb^{-1}) for masses up to about $(125 - 130) \text{ GeV}/c^2$. This prospect is the most powerful incentive we have for Run 2^{bis} . Over the range of masses accessible in associated production at the Tevatron, it should be possible to determine the mass of the Higgs boson to $\pm(1 - 3) \text{ GeV}/c^2$.

Recent studies [24] suggest that it may be possible to extend the reach of the Tevatron significantly by making use of the real- W -virtual- W (WW^*) decay modes for Higgs boson produced in the elementary reaction $gg \rightarrow H$. The WW^* channel has the largest branching fraction for $M_H \gtrsim 140 \text{ GeV}/c^2$. According to the analysis summarized in Figure 8, the large cross section \times branching fraction of the $gg \rightarrow H \rightarrow WW^*$ mode extends the $3\text{-}\sigma$ detection sensitivity of Run 2^{bis} into the region $145 \text{ GeV}/c^2 \lesssim M_H \lesssim 180 \text{ GeV}/c^2$. This is an extremely exciting opportunity, and it is important that the WW^* proposal receive independent critical analysis.

Constraints on the Higgs-Boson Mass

The electroweak theory itself provides reason to expect that discoveries will not end with the Higgs boson. Outside a narrow window of Higgs-boson masses, the electroweak theory cannot be complete.¹¹

Scalar field theories make sense on all energy scales only if they are noninteracting, or “trivial.” For any given Higgs-boson mass, there is a maximum energy scale Λ^* at which the theory ceases to make sense. Equivalently, if the theory is to be valid up to a certain scale Λ^* , that implies an upper bound on the Higgs-boson mass. The description of the Higgs boson as an elementary scalar is at best an effective theory, valid over a finite range of energies. A perturbative analysis identifies

$$\Lambda^* \leq M_H \exp\left(\frac{4\pi^2 v^2}{3M_H^2}\right), \quad (6)$$

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

This perturbative analysis breaks down when the Higgs-boson mass approaches 1 TeV/ c^2 and the interactions become strong. Lattice analyses indicate that, for the theory to describe physics to an accuracy of a few percent up to a few TeV, the mass of the Higgs boson can be no more than about 710 ± 60 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.

A lower bound is obtained by computing the first quantum corrections to the classical potential $V(\varphi^\dagger\varphi) = \mu^2(\varphi^\dagger\varphi) + |\lambda|(\varphi^\dagger\varphi)^2$. Requiring that $\langle\varphi\rangle_0 \neq 0$ be an absolute minimum of the one-loop potential up to a scale Λ yields the vacuum-stability condition

$$M_H^2 > \frac{3G_F\sqrt{2}}{8\pi^2}(2M_W^4 + M_Z^4 - 4m_t^4) \log(\Lambda^2/v^2). \quad (7)$$

The upper and lower bounds plotted in Figure 9 are the results of full two-loop calculations [25]. There I have also indicated the upper bound on M_H derived from precision electroweak measurements *in the framework of the standard electroweak theory*. 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 $\Lambda^* = 10^{16}$ GeV, then the Higgs-boson mass must lie in the interval $145 \text{ GeV}/c^2 \lesssim M_W \lesssim 170 \text{ GeV}/c^2$.

¹¹⁾ The origin of the mass bounds is reviewed in Ref. [15], where complete references will be found.

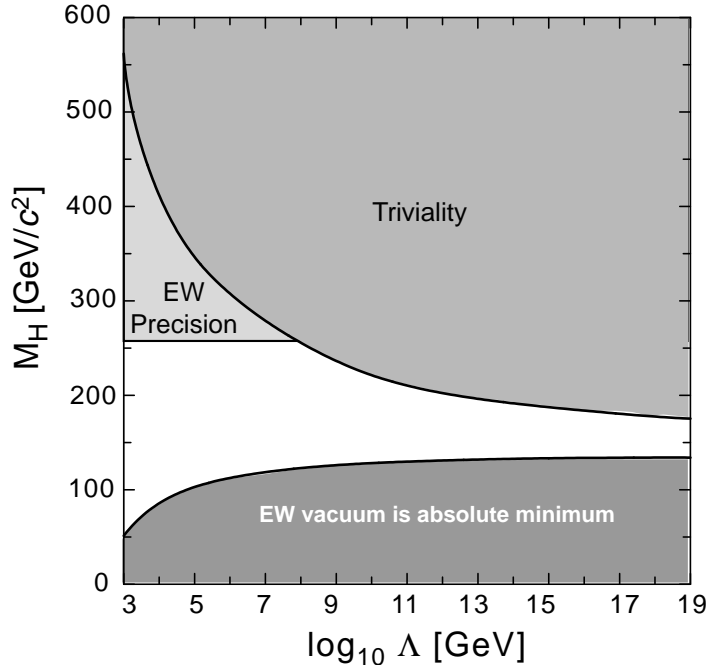


FIGURE 9. Bounds on the Higgs-boson mass that follow from requirements that the electroweak theory be consistent up to the energy scale Λ . The upper bound follows from triviality conditions; the lower bound follows from the requirement that the minimum of V occur for $\varphi \neq 0$. Also shown is the range of masses excluded at the 95% confidence level by precision measurements.

The Minimal Supersymmetric Standard Model

One of the best phenomenological motivations for supersymmetry on the 1-TeV scale is that the minimal supersymmetric extension of the standard model so closely approximates the standard model itself. A nice illustration of the small differences between predictions of supersymmetric models and the standard model is the compilation of pulls prepared by Erler and Pierce [26], which is shown in Figure 10. This is a nontrivial property of new physics beyond the standard model, and a requirement urged on us by the unbroken quantitative success of the established theory. On the aesthetic side, supersymmetry is the maximal—indeed, unique—extension of Poincaré invariance. It also offers a path to the incorporation of gravity, since local supersymmetry leads directly to supergravity. As a practical matter, supersymmetry on the 1-TeV scale offers a solution to the naturalness problem, and allows a fundamental scalar to exist at low energies.

Many recent papers investigate the constraints that precision electroweak measurements might place on the spectrum of superpartners. As a representative example, let me point to the analysis of Cho and Hagiwara [20] already cited for the standard-model fit. Their best overall fit, including the particles of the minimal supersymmetric standard model, suggests that the lightest readily observable superpartners should be (wino-like) charginos, with masses in the neighborhood

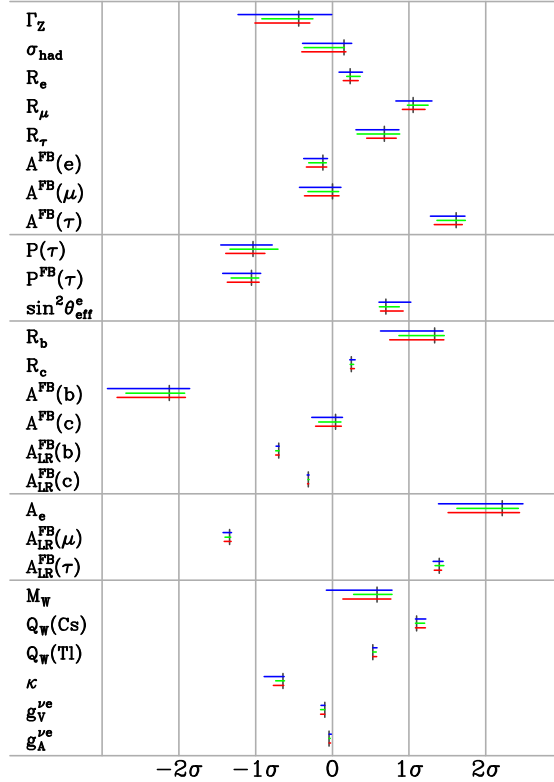


FIGURE 10. The range of best fit predictions of precision observables in the supergravity model (upper horizontal lines), the $\mathbf{5} \oplus \mathbf{5}^*$ gauge-mediated model (middle lines), the $\mathbf{10} \oplus \mathbf{10}^*$ gauge-mediated model (lower lines), and in the standard model at its global best fit value (vertical lines), in units of standard deviation, from Ref. [26].

of $100 \text{ GeV}/c^2$. The mass of the lightest neutralino, $\tilde{\chi}_1^0$, is about $50 \text{ GeV}/c^2$, while most of the doublet squarks weigh more than about $200 \text{ GeV}/c^2$. The sleptons are relatively heavy, the lightest being the $\tilde{\tau}_1$, in the range $134 - 169 \text{ GeV}/c^2$. For the moment, I think it is prudent to take statements of this kind as mildly suggestive. As we learn more about the mass of the Higgs boson, and constrain the standard-model contributions to electroweak observables still further, our ability to derive definite expectations for physics beyond the standard model should grow.

Because the minimal supersymmetric standard model (MSSM) implies *upper bounds* on the mass of the lightest scalar h^0 , it sets attractive targets for experiment. Two such upper bounds are shown as functions of the top-quark mass in Figure 11. The large- $\tan \beta$ limit of a general MSSM yields the upper curve; an infrared-fixed-point scheme with b - τ unification produces an upper bound characterized by the lower curve. The vertical band shows the current information on m_t . The projected sensitivity of LEP 2 experiments [28] covers the full range of lightest-Higgs masses that occur in the infrared-fixed-point scheme. The sensitivity promised by Run 2^{bis} of the Tevatron [23] gives full coverage of h^0 masses in the MSSM. These are very intriguing experimental possibilities. For further discussion, consult the LEP 2

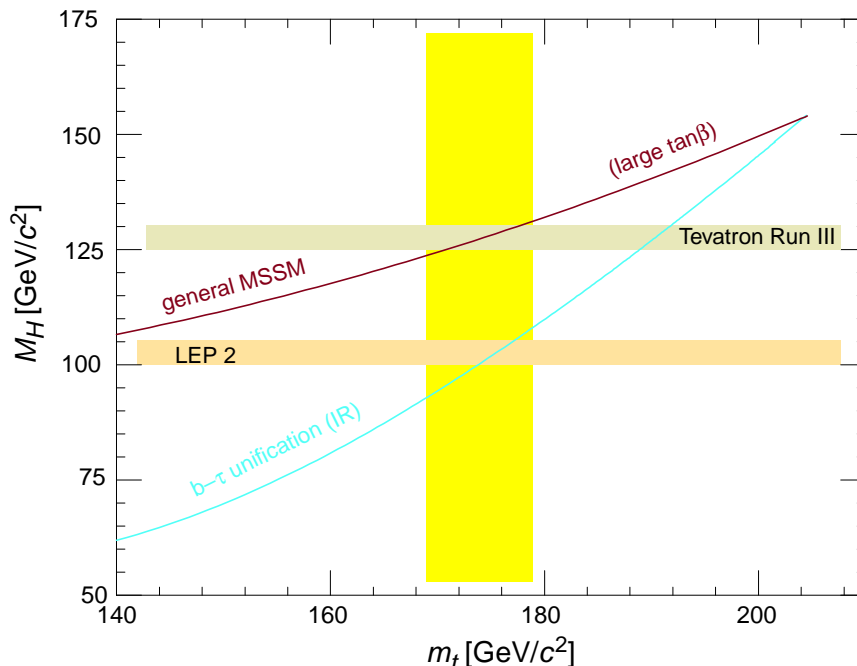


FIGURE 11. Upper bounds on the mass of the lightest Higgs boson, as a function of the top-quark mass, in two variants of the minimal supersymmetric standard model. The upper curve refers to a general MSSM, in the large- $\tan\beta$ limit; the lower curve corresponds to an infrared-fixed-point scenario with b - τ unification, from Ref. [27].

Yellow Book [29] and the forthcoming Proceedings of the Tevatron Supersymmetry / Higgs Workshop.¹²

How Much Can We Trust the Standard-Model Analysis of Precision Electroweak Data to Bound M_H ?

As suggestive as we may find the standard-model fits to precision electroweak data that tell us the Higgs boson is light, we know that the standard model must break down at a nearby energy scale if the Higgs boson is indeed light. We must ask, with some urgency, how the occurrence of new physics near the 1-TeV scale would modify our expectations for M_H . Hall and Kolda [30] have given one provocative illustration. When they include higher-dimension operators that could arise if the dimensionality of spacetime is greater than $3 \oplus 1$, their fit to the electroweak observables leads to no constraint on the Higgs-boson mass: *the constraints of the standard-model analysis simply evaporate*. That is a dramatic result, but we do not know whether the situation they envisage actually arises in a real theory, and whether the vacuum of that theory indeed breaks the electroweak symmetry. In a

¹²⁾ Current drafts are available at <http://fnth37.fnal.gov/susy.html>.

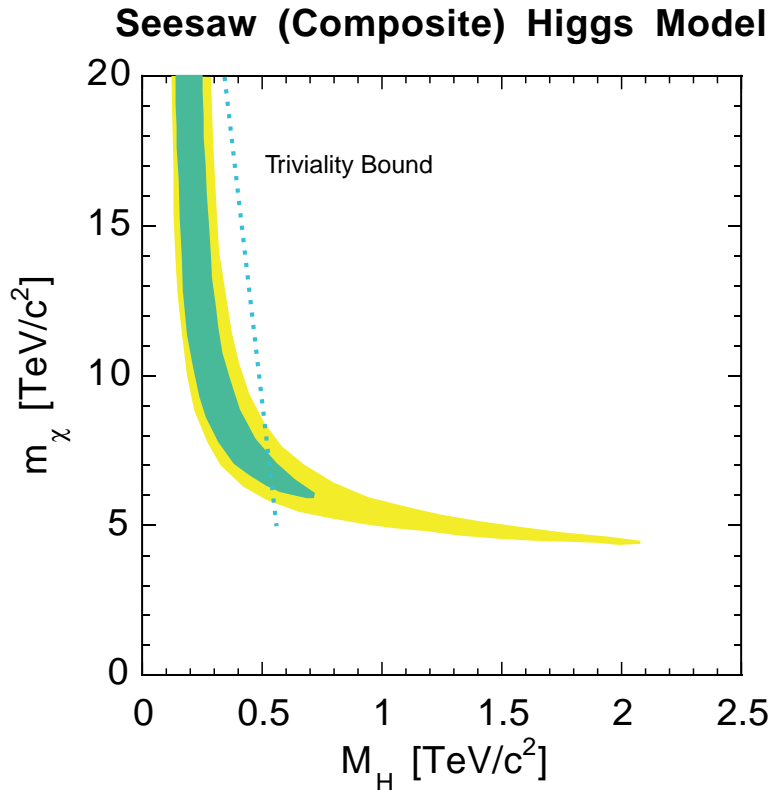


FIGURE 12. Precision electroweak constraints (at 68% and 90% CL) on the Higgs-boson mass and heavy-fermion mass in a top-quark seesaw model of a composite Higgs boson, [32]. The triviality bound is from the work of Chivukula and Evans, Ref. [33].

similar spirit, Rizzo and Wells [31] have shown that if the Higgs boson is trapped on a $3 \oplus 1$ -dimensional wall with the fermions, large Higgs masses (up to $500 \text{ GeV}/c^2$) and relatively light Kaluza-Klein mass scales provide a good fit to precision data.

More apposite, because it does correspond to a fully elaborated theory, is the topcolor seesaw example recently offered by Collins, Grant, and Georgi [32]. In a model that includes one additional heavy weak-singlet fermion χ , with weak hypercharge $Y(\chi) = \frac{4}{3}$, they find that the mass of the Higgs boson can exceed $300 \text{ GeV}/c^2$ (the $\gtrsim 90\%$ CL upper bound from standard-model analyses) for a heavy-fermion mass in the range $5 \text{ TeV}/c^2 \lesssim m_\chi \lesssim 7 \text{ TeV}/c^2$. The $1\text{-}\sigma$ and 90% CL allowed regions in the (m_χ, M_H) plane are shown in Figure 12.¹³

We need more examples based on plausible—or at least provocative—extensions of the standard model to help us understand how seriously to take the standard-model clues, and to help us think about what it might mean if the standard-model hopes are dashed.

¹³⁾ A second topcolor seesaw model based on two heavy fermions and two Higgs doublets leads to heavy scalars h^0, H^0, H^\pm , but a *light pseudoscalar*—quite different from the spin-0 spectrum of the MSSM.

Looking for Trouble

Suppose, in the face of the spectacular successes of the electroweak theory, we go looking for trouble. Where might we find it? The great mass of the top quark gives rise to the theoretical suspicion that anomalies are most likely to show themselves in the third generation of quarks and leptons. As it happens, the largest pull in precision measurements on the Z^0 pole involves b quarks. The forward-backward asymmetry for $b\bar{b}$ events measured at LEP and the left-right forward-backward asymmetry for $b\bar{b}$ events measured at SLD indicate a three-standard-deviation difference from the standard model for

$$A_b = \frac{L_b^2 - R_b^2}{L_b^2 + R_b^2}, \quad (8)$$

where L_b and R_b are the left-handed and right-handed chiral couplings of the Z to b quarks. Pursuing this line, we arrive at the suggestion that $\delta R_b/R_b^{\text{theory}} \approx -40\%$, whereas $\delta L_b/L_b^{\text{theory}} \approx 1\%$. Therefore, if we want to find deviations from standard-model predictions, we should try to isolate effects of the right-handed b coupling [34]. If this anomaly is real, we might expect to observe flavor-changing neutral-current transitions $b \rightarrow s$, $b \rightarrow d$, and $s \rightarrow d$.

Bennett and Wieman (Boulder) have reported a new determination of the weak charge of Cesium by measuring the transition polarizability for the 6S-7S transition [35]. The new value,

$$Q_W(\text{Cs}) = -72.06 \pm 0.28 \text{ (expt)} \pm 0.34 \text{ (theory)}, \quad (9)$$

represents a sevenfold improvement in the experimental error and a significant reduction in the theoretical uncertainty. It lies about 2.5 standard deviations above the prediction of the standard model. We are left with the traditional situation in which elegant measurements of parity nonconservation in atoms are on the edge of incompatibility with the standard model.

A number of authors [21] have noted that the discrepancy in the weak charge Q_W and a $2\text{-}\sigma$ anomaly¹⁴ in the total width of the Z^0 can be reduced by introducing a Z' boson with a mass of about $800 \text{ GeV}/c^2$. The additional neutral gauge boson resembles the Z_χ familiar from unified theories based on the group E_6 .

The Vacuum Energy Problem

I want to spend a moment to revisit a longstanding, but usually unspoken, challenge to the completeness of the electroweak theory as we have defined it: the vacuum energy problem [36]. I do so not only for its intrinsic interest, but also to raise the question, "Which problems of completeness and consistency do we worry

¹⁴) In Erler & Langacker's fit, for example, the number of light neutrino species inferred from the invisible width of the Z^0 is $N_\nu = 2.985 \pm 0.008 = 3 - 2\sigma$.

about at a given moment?” It is perfectly acceptable science—indeed, it is often essential—to put certain problems aside, in the expectation that we will return to them at the right moment. What is important is never to forget that the problems are there, even if we do not allow them to paralyze us.

For the usual Higgs potential, $V(\varphi^\dagger\varphi) = \mu^2(\varphi^\dagger\varphi) + |\lambda|(\varphi^\dagger\varphi)^2$, the value of the potential at the minimum is

$$V(\langle\varphi^\dagger\varphi\rangle_0) = \frac{\mu^2 v^2}{4} = -\frac{|\lambda| v^4}{4} < 0. \quad (10)$$

Identifying $M_H^2 = -2\mu^2$, we see that the Higgs potential contributes a field-independent constant term,

$$\varrho_H \equiv \frac{M_H^2 v^2}{8}. \quad (11)$$

I have chosen the notation ϱ_H because the constant term in the Lagrangian plays the role of a vacuum energy density. When we consider gravitation, adding a vacuum energy density ϱ_{vac} is equivalent to adding a cosmological constant term to Einstein’s equation. Although recent observations¹⁵ raise the intriguing possibility that the cosmological constant may be different from zero, the essential observational fact is that the vacuum energy density must be very tiny indeed,¹⁶

$$\varrho_{\text{vac}} \lesssim 10^{-46} \text{ GeV}^4. \quad (12)$$

Therein lies the puzzle: if we take $v = (G_F\sqrt{2})^{-\frac{1}{2}} \approx 246 \text{ GeV}$ and insert the current experimental lower bound [28] $M_H \gtrsim 105 \text{ GeV}/c^2$ into (11), we find that the contribution of the Higgs field to the vacuum energy density is

$$\varrho_H \gtrsim 8 \times 10^7 \text{ GeV}^4, \quad (13)$$

some 54 orders of magnitude larger than the upper bound inferred from the cosmological constant.

What are we to make of this mismatch? The fact that $\varrho_H \gg \varrho_{\text{vac}}$ means that the smallness of the cosmological constant needs to be explained. In a unified theory of the strong, weak, and electromagnetic interactions, other (heavy!) Higgs fields have nonzero vacuum expectation values that may give rise to still greater mismatches. At a fundamental level, we can therefore conclude that a spontaneously broken gauge theory of the strong, weak, and electromagnetic interactions—or merely of the electroweak interactions—cannot be complete. Either we must find a separate principle to zero the vacuum energy density of the Higgs field, or we may suppose that a proper quantum theory of gravity, in combination with the other interactions, will resolve the puzzle of the cosmological constant. The vacuum energy problem must be an important clue. But to what?

¹⁵⁾ For a cogent summary of current knowledge of the cosmological parameters, including evidence for a cosmological constant, see Ref. [37].

¹⁶⁾ For a useful summary of gravitational theory, see the essay by T. d’Amour in §14 of the 1998 *Review of Particle Physics*, Ref. [11].

THE PROBLEMS OF MASS, AND OF MASS SCALES

Electroweak symmetry breaking sets the values of the W - and Z -boson masses. At tree level in the electroweak theory, we have

$$M_W^2 = g^2 v^2 / 2 = \pi \alpha / G_F \sqrt{2} \sin^2 \theta_W, \quad (14)$$

$$M_Z^2 = M_W^2 / \cos^2 \theta_W, \quad (15)$$

where the electroweak scale is $v = (G_F \sqrt{2})^{-\frac{1}{2}} \approx 246$ GeV. But the electroweak scale is not the only scale. It seems certain that we must also consider the Planck scale, derived from the strength of Newton's constant,

$$M_{\text{Planck}} = (\hbar c / G_{\text{Newton}})^{\frac{1}{2}} \approx 1.22 \times 10^{19} \text{ GeV}. \quad (16)$$

It is also probable that we must take account of the $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ unification scale around 10^{15-16} GeV, and there may well be a distinct flavor scale. The existence of other scales is behind the famous problem of the Higgs scalar mass: how to keep the distant scales from mixing in the face of quantum corrections, or how to stabilize the mass of the Higgs boson on the electroweak scale.

It is because G_{Newton} is so small (or because M_{Planck} is so large) that we normally consider gravitation irrelevant for particle physics. The graviton-quark-antiquark coupling is generically $\sim E/M_{\text{Planck}}$, so it is easy to make a dimensional estimate of the branching fraction for a gravitationally mediated rare kaon decay: $B(K_L \rightarrow \pi^0 G) \sim (M_K/M_{\text{Planck}})^2 \sim 10^{-38}$, which is truly negligible!

We know from the electroweak theory alone that the 1-TeV scale is special. Partial-wave unitarity applied to gauge-boson scattering tells us that unless the Higgs-boson mass respects

$$M_H^2 < \frac{8\pi\sqrt{2}}{3G_F} \approx (1 \text{ TeV}/c^2)^2, \quad (17)$$

new physics is to be found on the 1-TeV scale [38]. To stabilize the Higgs-boson mass against uncontrolled quantum corrections, and to resolve the mass-hierarchy problem, we consider electroweak physics beyond the standard model. The most promising approaches are to generalize $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ to a theory with a composite Higgs boson in which the electroweak symmetry is broken dynamically (technicolor and related theories) or to a supersymmetric standard model.

Let us look a little further at the problem of fermion masses.¹⁷ In the electroweak theory, the value of each quark or charged-lepton mass is set by a new, unknown, Yukawa coupling. We define the left-handed doublets and right-handed singlets

$$\underline{L}_e = \begin{pmatrix} \nu_e \\ e \end{pmatrix}_L, \quad \underline{R}_e \equiv e_R; \quad \underline{L}_q = \begin{pmatrix} u \\ d \end{pmatrix}_L, \quad \underline{R}_u \equiv u_R, \quad \underline{R}_d \equiv d_R. \quad (18)$$

¹⁷⁾ For an overview of the standard-model approach to fermion mass, see Ref. [39].

Then the electron's Yukawa term in the electroweak Lagrangian is

$$\mathcal{L}_{\text{Yukawa}}^{(e)} = -\zeta_e [\bar{\mathbf{R}}_e (\varphi^\dagger \mathbf{L}_e) + (\bar{\mathbf{L}}_e \varphi) \mathbf{R}_e] , \quad (19)$$

where φ is the Higgs field, so that the electron mass is $m_e = \zeta_e v / \sqrt{2}$. Similar expressions obtain for the quark Yukawa couplings. Inasmuch as we do not know how to calculate the fermion Yukawa couplings ζ_f , I believe that *we should consider the sources of all fermion masses as physics beyond the standard model*.

The values of the Yukawa couplings are vastly different for different fermions: for the top quark, $\zeta_t \approx 1$, for the electron $\zeta_e \approx 3 \times 10^{-6}$, and if the neutrinos have Dirac masses, presumably $\zeta_\nu \approx 10^{-10}$.¹⁸ What accounts for the range and values of the Yukawa couplings? Our best hope until now has been the suggestion from unified theories that the pattern of fermion masses simplifies on high scales. The classic intriguing prediction of the $SU(5)$ unified theory involves the masses of the b quark and the τ lepton, which are degenerate at the unification point for a simple pattern of spontaneous symmetry breaking. The different running of the quark and lepton masses to low scales then leads to the prediction $m_b \approx 3m_\tau$, in suggestive agreement with what we know from experiment.

The conventional approach to new physics has been to extend the standard model to understand why the electroweak scale (and the mass of the Higgs boson) is so much smaller than the Planck scale. A novel approach that has been developed over the past two years is instead to *change gravity* to understand why the Planck scale is so much greater than the electroweak scale [40]. Now, experiment tells us that gravitation closely follows the Newtonian force law down to distances on the order of 1 mm. Let us parameterize deviations from a $1/r$ gravitational potential in terms of a relative strength ε_G and a range λ_G , so that

$$V(r) = - \int dr_1 \int dr_2 \frac{G_{\text{Newton}} \rho(r_1) \rho(r_2)}{r_{12}} [1 + \varepsilon_G \exp(-r_{12}/\lambda_G)] , \quad (20)$$

where $\rho(r_i)$ is the mass density of object i and r_{12} is the separation between bodies 1 and 2. Elegant experiments that study details of Casimir and Van der Waals forces imply bounds on anomalous gravitational interactions, as shown in Figure 13. Below about a millimeter, the constraints on deviations from Newton's inverse-square force law deteriorate rapidly, so nothing prevents us from considering changes to gravity even on a small but macroscopic scale.

For its internal consistency, string theory requires an additional six or seven space dimensions, beyond the $3 + 1$ dimensions of everyday experience. Until recently it has been presumed that the extra dimensions must be compactified on the Planck scale, with a compactification radius $R_{\text{unobserved}} \approx 1/M_{\text{Planck}} \approx 1.6 \times 10^{-35}$ m. The new wrinkle is to consider that the $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ standard-model gauge fields, plus needed extensions, reside on $3 + 1$ -dimensional branes, not in the extra dimensions, but that gravity can propagate into the extra dimensions.

¹⁸⁾ I am quoting the values of the Yukawa couplings at a low scale typical of the masses themselves.

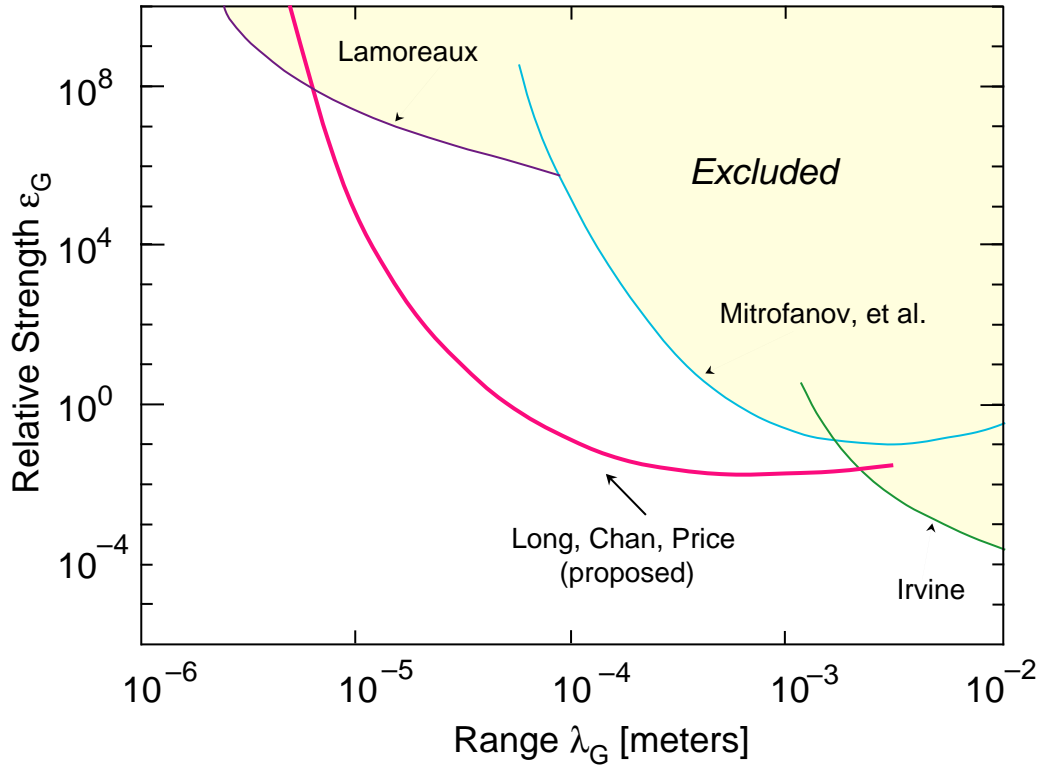


FIGURE 13. Experimental limits on the strength ε_G (relative to gravity) versus the range λ_G of a new long-range force, together with the anticipated sensitivity of a new experiment based on small mechanical resonators [41].

How does this hypothesis change the picture? The dimensional analysis (Gauss's law, if you like) that relates Newton's constant to the Planck scale changes. If gravity propagates in n extra dimensions with radius R , then

$$G_{\text{Newton}} \sim M_{\text{Planck}}^{-2} \sim M^*{}^{-n-2} R^{-n}, \quad (21)$$

where M^* is gravity's true scale. Notice that if we boldly take M^* to be as small as $1 \text{ TeV}/c^2$ (which other constraints may not allow), then the radius of the extra dimensions is required to be smaller than about 1 mm , for $n \geq 2$. If we use the four-dimensional force law to extrapolate the strength of gravity from low energies to high, we find that gravity becomes as strong as the other forces on the Planck scale, as shown by the dashed line in Figure 14. If the force law changes at an energy $1/R$, as the large-extra-dimensions scenario suggests, then the forces are unified at an energy M^* , as shown by the solid line in Figure 14. What we know as the Planck scale is then a mirage that results from a false extrapolation: treating gravity as four-dimensional down to arbitrarily small distances, when in fact—or at least in this particular fiction—gravity propagates in $3 + n$ spatial dimensions. The Planck mass is an artifact, given by $M_{\text{Planck}} = M^*(M^*R)^{n/2}$.

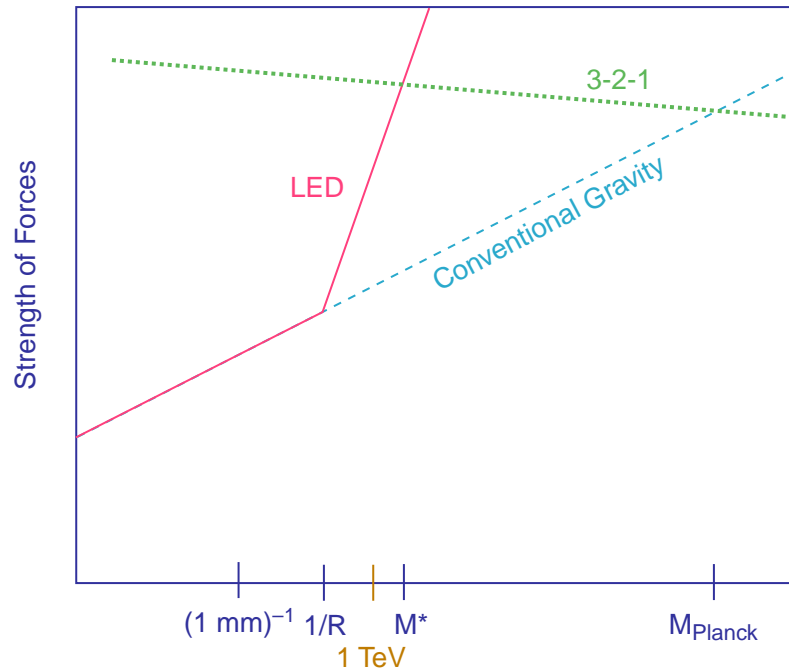


FIGURE 14. One of these extrapolations (at least!) is false.

Although the idea that extra dimensions are just around the corner—either on the submillimeter scale or on the TeV scale—is preposterous, it is not ruled out by observations. For that reason alone, we should entertain ourselves by entertaining the consequences. Many authors have considered the gravitational excitation of a tower of Kaluza–Klein modes in the extra dimensions, which would give rise to a missing (transverse) energy signature in collider experiments [42]. We call these excitations *provatons*, after the Greek word for a sheep in a flock.

“Large” extra dimensions present us with new ways to think about the exponential range of Yukawa couplings. Arkani-Hamed, Schmaltz, and Mirabelli [43] have speculated that if the standard-model brane has a small thickness, the wave packets representing different fermion species might have different locations within the extra dimension, as depicted in Figure 15. On this picture, the Yukawa couplings measure the overlap in the extra dimensions of the left-handed and right-handed wave packets and the Higgs field, presumed pervasive. Exponentially large differences might then arise from small offsets in the new coordinate(s). True or not, it is a completely different way of looking at an important problem. Speculations of this sort give a new force to the useful metaphor of particle accelerators as giant microscopes. We may soon be able to examine Nature on such a small scale that we uncover not only fine new features of the familiar ground, but also reveal hitherto unknown dimensions.

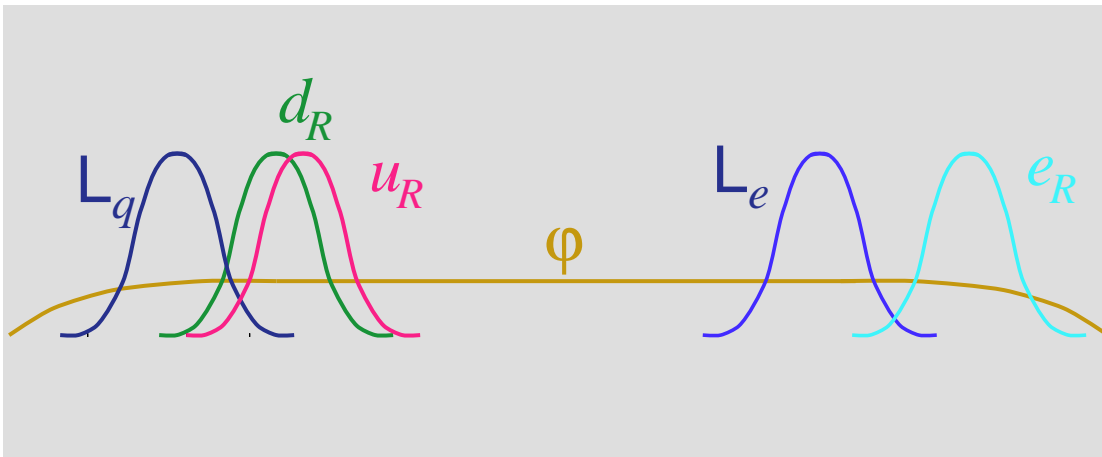


FIGURE 15. The locations in the brane of wave packets that represent different fermion species.

QCD AND THE ELECTROWEAK THEORY

Because QCD is asymptotically free and becomes strong at low energies, it has a rich phase structure. It is worth asking whether the phases that arise in QCD hold lessons for electroweak symmetry breaking. We already know of one case for which the answer is Yes.

Our conception of electroweak symmetry breaking is modeled upon our understanding of the superconducting phase transition. The Higgs mechanism of the standard model is the relativistic generalization of the Ginzburg-Landau description of the superconducting phase transition. The macroscopic order parameter of the Ginzburg-Landau phenomenology corresponds to the wave function of superconducting charge carriers, which acquires a nonzero vacuum expectation value in the superconducting state and generates a photon mass that explains the Meissner effect. The microscopic Bardeen-Cooper-Schrieffer theory identifies the dynamical origin of the order parameter with the formation of correlated states of elementary fermions, the Cooper pairs of electrons. Can we find a dynamical mechanism for electroweak symmetry breaking that corresponds to the BCS theory?

The elementary fermions—electrons—and 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 $\sim \Lambda_{\text{QCD}}$, the strong interactions become strong, fermion condensates appear, and the chiral symmetry is spontaneously broken to the familiar flavor (isospin) symmetry:

$$SU(2)_L \otimes SU(2)_R \rightarrow SU(2)_V . \quad (22)$$

Three Goldstone bosons appear, one for each broken generator of the original chiral invariance. These were identified by Nambu [44] as three massless pions.

The broken generators are three axial currents whose couplings to pions are measured by the pion decay constant $f_\pi = 93$ MeV. 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 $\sim gf_\pi$. The mass-squared matrix,

$$\mathcal{M}^2 = \begin{pmatrix} g^2 & 0 & 0 & 0 \\ 0 & g^2 & 0 & 0 \\ 0 & 0 & g^2 & gg' \\ 0 & 0 & gg' & g'^2 \end{pmatrix} \frac{f_\pi^2}{4}, \quad (23)$$

(where the rows and columns correspond to the weak-isospin gauge bosons W^+ , W^- , W_3 , and the weak-hypercharge gauge boson \mathcal{A}) has the same structure as the mass-squared matrix for gauge bosons in the standard electroweak theory. Diagonalizing the matrix (23), we find that $M_W^2 = g^2 f_\pi^2/4$ and $M_Z^2 = (g^2 + g'^2) f_\pi^2/4$, so that

$$\frac{M_Z^2}{M_W^2} = \frac{(g^2 + g'^2)}{g^2} = \frac{1}{\cos^2 \theta_W}. \quad (24)$$

The photon emerges massless.

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; $M_W \approx 30$ MeV/ c^2 and $M_Z \approx 34$ MeV/ c^2 , about $\frac{1}{2650} \times$ the true masses [45].

If only we didn't know f_π ! The idea of replacing the elementary Higgs boson with a fermion-antifermion condensate is too good to abandon without a struggle [46]. Perhaps a more general formulation of the chiral-symmetry-breaking idea has merit. We replace the massless up and down quarks with new “technifermions” and replace QCD with a new “technicolor” gauge interaction [47,48]. We choose the scale of the interaction—the analogue of Λ_{QCD} —so that

$$f_\pi \rightarrow F_\pi = v = (G_F \sqrt{2})^{-1/2}. \quad (25)$$

Repeating the analysis we have just made for QCD, we predict the correct (tree-level) masses for W^\pm , Z^0 , and the photon. 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.

Technicolor shows how the generation of intermediate boson masses could arise from strong dynamics (though not the strong dynamics of QCD), 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 [49]. Consequently, technicolor serves as a reminder that there are two problems of mass: explaining the masses of the gauge bosons, which demands an understanding of electroweak symmetry breaking; and accounting for the quark and lepton masses, which requires not only an understanding of electroweak symmetry breaking but also a theory of the Yukawa couplings that set the scale of fermion masses in the standard model. We can be confident that the origin of gauge-boson masses will be understood on the 1-TeV scale. We do not know where we will decode the pattern of the Yukawa couplings.

Is it possible that other interesting phases of QCD—color superconductivity [50,51], for example—might hold lessons for electroweak symmetry breaking under normal or unusual conditions?

CONCLUDING REMARKS

Wonderful opportunities await particle physics over the next decade, with the coming of the Large Hadron Collider at CERN to explore the 1-TeV scale (extending the efforts at LEP and the Tevatron to unravel the nature of electroweak symmetry breaking) and many initiatives to develop our understanding of the problem of identity: what makes a neutrino a neutrino and a top quark a top quark. Here I have in mind the work of the B factories and the Tevatron collider on CP violation and the weak interactions of the b quark; the wonderfully sensitive experiments at Brookhaven, Fermilab, CERN, and Frascati on CP violation and rare decays of kaons; the prospect of definitive accelerator experiments on neutrino oscillations and the nature of the neutrinos; and a host of new experiments on the sensitivity frontier. We might even learn to read experiment for clues about the dimensionality of spacetime. If we are inventive enough, we may be able to follow this rich menu with the physics opportunities offered by a linear collider and a (muon storage ring) neutrino factory. I expect a remarkable flowering of experimental particle physics, and of theoretical physics that engages with experiment.

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