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Substructure and Strong Interactions at the TeV Scale

MICHAEL E. PESKIN*

Stanford Linear Accelerator Center
Stanford University, Stanford, California, 94305

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SUBSTRUCTURE AND STRONG INTERACTIONS AT THE TEV SCALE

Michael E. Peskin

Stanford Linear Accelerator Center
Stanford University, Stanford, California, 94305

This is the first of two lectures at this Symposium devoted to possible modifications of the standard model of strong, weak, and electromagnetic interactions in the 100 GeV – 1000 GeV energy range. Let me begin by explaining why it is important to contemplate the breakdown of our current theoretical picture. The reason is most certainly not that the standard model—the gauge theory of $SU(3)$ color $\times$ weak interaction $SU(2) \times U(1)$—has in some way failed a crucial experimental test. Indeed, almost every talk at this conference has offered a new, nontrivial confirmation of $SU(3) \times SU(2) \times U(1)$. Rather, we have reached the stage where the most pressing problems in elementary particle physics are problems which the standard model is not equipped to answer. Outstanding among these are, first, the question of the origin of the W and Z boson masses and the breaking of the weak-interaction gauge symmetry, and second, the problem of the origin of the quark and lepton masses. The standard model addresses these two questions only by providing adjustable parameters which account for these effects. The prediction of the values of these parameters lies outside the realm of the model, in exactly the same sense that the prediction of the value of the fine structure constant lies outside the realm of pure quantum electrodynamics.

The standard model does, however, provide one important insight into the nature of these effects: It insists that they are tightly connected to one another. One can see this in either of two ways. First, the standard model assigns to left- and right-handed fermions different quantum numbers. A fermion mass term, which is essentially a coupling of the right- and left-handed fermion species, can appear only when the weak gauge symmetry $SU(2) \times U(1)$ is broken. Second, the mass scales of gauge symmetry breaking and fermion masses are relatively close to one another, at least when one considers the grand sweep of scales contemplated in grand unified theories. One often hears the question of what drives the weak symmetry breaking phrased as a *gauge hierarchy problem*, the question of why the distance between the weak scale and some more fundamental scale such as the Planck mass or the scale of grand unification is so large. The standard model links to this hierarchy the hierarchy of masses which yields, eventually, the masses of the fermions most familiar to our experience.

If we are to seek a joint solution to these problems, we must begin in seeking the origin of the weak interaction mass scale. In the simplest version of the standard model, this scale is characterized by the expectation value of the fundamental Higgs field:

$$\langle \phi \rangle = 240 \text{ GeV} . \tag{1}$$

To explain this scale, we must replace the fundamental Higgs, which is essentially just a parametrization, by some comprehensible mechanics. The general magnitude of the energies involved in this mechanism will be that of (1), roughly 1 TeV. Such explanations fall naturally into two classes, those in which this mechanism is essentially perturbative, and those in which nonperturbative physics plays an essential role. The most successful of the perturbative approaches has been that of supersymmetry. John Ellis will discuss the status of this approach in great detail in his lecture at this Symposium. My lecture will review models which make use of nonperturbative physics, models which require new bound states and new strong interactions in the energy region near 1 TeV.

The outline of this lecture is as follows: I will first give a review of the current status of the three main theoretical ideas relevant to strong-interaction 1 TeV physics. All involve the assumption that some object which is assumed to be fundamental in the standard model actually has dynamical internal structure. I will discuss, in turn, the ideas of composite vector bosons, Higgs bosons ("Technicolor"), and matter fermions. In general, I will be discussing complex, mechanistic models of the new physics, models which are not especially beautiful but which have, potentially, the power to explain. I will then enter a brief digression on how the weak interaction allow us to probe for this new structure. Finally, I will discuss direct manifestations of new 1 TeV strong interactions. Remarkably, as one actually reaches the TeV energy scale experimentally, the consequences of new strong interactions become as dramatic as the specific models of these interactions are obscure.

Theoretical Ideas about Compositeness

In this section, I will review ideas concerning the possible compositeness of the various components of the standard model and the experimental tests which constrain those ideas. In general, I will be discussing classes of models rather than specific schemes. Some of the theoretical constructions will only be described only sketchily; a more complete, though still introductory, theoretical discussion can be found in my lecture at the 1981 Symposium.[1]

In discussing the experimental constraints on compositeness, I would like to make statements as model-independent as possible. A useful tool for identifying such model-independent information is the method of *effective Lagrangians*, pioneered by Wilson[2] in the 1960's and shown by Weinberg[3] to be a method of great power when applied to gauge theories. In a theory with a regime of weak coupling, such as we find in the standard model, it is essentially enlightened dimensional analysis. The method consists of writing the most general Lagrangian consistent with gauge and global symmetries which remain unbroken at a given scale, and then fixing from dimensional analysis and a general picture of the symmetry-breaking pattern the magnitude of the coefficients in this Lagrangian. Consider, for example, the general theory of quarks:

$$\mathcal{L}_{\text{eff}} \;=\; \bar{q} i \not{D} q \;+\; \lambda \bar{q}_L \phi q_R \;+\; \cdots \;+\; \frac{1}{\Lambda^2} \bar{q}_L \gamma^\mu q_L \bar{q}_R \gamma_\mu q_R \;+ \cdots\,, \qquad (2)$$

where $q_{R,L} = \frac{1}{2}(1 \pm \gamma^5) q$. The quark fields are normalized so that the first term has coefficient 1; this fixes all of the gauge couplings of the quarks. The second term is the simplest one which can flip quark helicity and is allowed by the $SU(2) \times U(1)$ weak-interaction symmetry. Its coefficient is dimensionless and should then be of order 1 unless it was somehow forbidden at a deeper level of the theory with a higher symmetry. It is the dynamics of this deeper level, then, which suppresses large quark mixing angles, and which makes the value of λ very small for particular quarks (*e.g.*, 10^{-5} for the u quark). The third term given is a 4-fermion term of fairly general structure. Its coefficient has the dimensions of $(\text{mass})^{-2}$. The coupling shown does not appear in the standard model, so we would expect the size of the mass Λ to reflect the scale at which the standard model yields receives corrections. I will assume that any term that can appear in an effective Lagrangian will appear, and that no term which arises from new strong interactions should be suppressed by powers of a small coupling constant.

• Composite Gauge Bosons?

I will first discuss the question of whether the weak vector bosons W and Z might be composite. In quantum field theory, general theorems prohibit the appearance of massless vector particles which carry nontrivial values of a conserved charge, unless both the bosons and the current of this charge are pieces of the structure of an exact gauge symmetry.[4–6] However, the W and Z bosons are not massless, and this has led many authors to speculate that they might, in fact, be dynamically generated. One often hears the statement[7] that, since all other short-range interactions (for example, interatomic forces and the nuclear potential) are built up from composite particle exchanges, the weak interactions should be built in the same way. It is important to look into this point critically as we enter the era in which the W and Z are produced directly, so that their properties can be studied with precision. Many of the properties already observed follow from the gauge structure of the standard model. We must ask to what extent they follow from less restrictive principles, and how one might make a test which could single out theories in which the W and Z are fundamental gauge bosons.

We should begin with a crucial observation of Bjorken[8] that the successes of the standard model in low-energy reactions follow from a principle much weaker than $SU(2) \times U(1)$ gauge invariance, the principle of an $SU(2)$ invariance of whatever new sector is responsible for the weak interactions. All successes of the standard model at energies well below the W mass can be summarized as supporting the effective current-current interaction

$$\delta \mathcal{L}_{\text{eff}} = \frac{G_F}{2}\left[\sqrt{2} J_L^{\mu +} J_L^{\mu -} + \left(J_L^{\mu 3} - \sin^2\theta_w J_{EM}^{\mu}\right)^2\right], \tag{3}$$

which is isospin-invariant up to the coupling to electromagnetism. In the standard model, this form follows from the fact that the Higgs structure which breaks the gauge symmetry preserves an $SU(2)$ isospin symmetry under which the three $SU(2)$ bosons transform as a triplet. This symmetry, called *custodial* $SU(2)$, plus the assumption that a photon must remain massless, forces the mass matrix for the $SU(2) \times U(1)$ bosons to take the form[9] :

$$m^2 = \frac{\langle\phi\rangle^2}{4}\begin{pmatrix} g^2 & & & \\ & g^2 & & \\ & & g^2 & gg' \\ & & gg' & g'^2 \end{pmatrix} \begin{pmatrix} W^1 \\ W^2 \\ W^3 \\ B \end{pmatrix} . \tag{4}$$

This mass matrix leads directly to (3), and also to a simple relation for the physical boson masses:[10]

$$\rho = \frac{m_W^2}{m_Z^2 \cos^2\theta_w} = 1 . \tag{5}$$

Gauge invariance plays a major role in determining the detailed structure of the weak interactions, however, at energies above the masses of the W and Z. The reason for this is that theories of interacting massive vector bosons are very susceptible to violation of unitarity; gauge invariance provides the only simple cure for this difficulty. To understand the problem, recall that a boson boosted to $k^\mu = (E, 0, 0, k)$ has its third polarization vector boosted to

$$\epsilon_L = (\frac{k}{m}, 0, 0, \frac{E}{m}) . \tag{6}$$

The diagram shown in Fig. 1(a), considered alone, yields a rate for the process $e^+e^- \rightarrow W^+W^-$, in units of R, of order

$$\left(\epsilon_L(1)\cdot\epsilon_L(2)\right)^2 \sim (\frac{s}{m_W^2})^2 . \tag{7}$$

In the standard model, however, one must sum the three diagrams shown in Fig. 1(b). The complete calculation contains dramatic, apparently miraculous, cancellations, and yields a rate of order 1.[11,12] Lee, Quigg, and Thacker[13] have studied the more hypothetical process $W^+W^- \rightarrow W^+W^-$ and have shown that, here, one requires also the W-W-Higgs coupling of the sum of lowest-order diagrams is to be consistent with unitarity. If the theory which contains the W bosons is not gauge-invariant, then, it must be only an effective theory in an energy range $\sqrt{s} << \Lambda < 1$ TeV in which the leading-order amplitudes are still unitary.

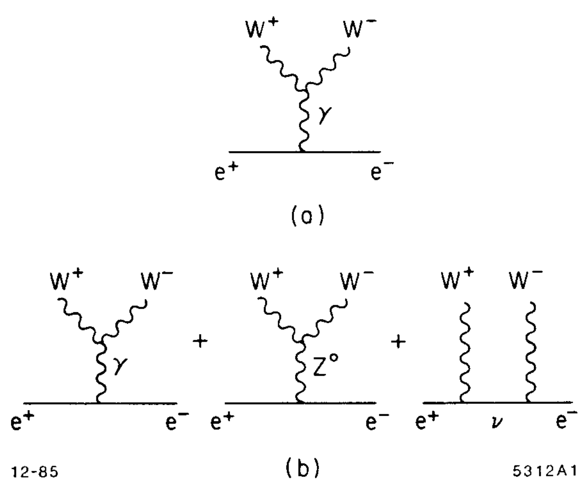


Fig. 1. (a) One diagram contributing to $e^+e^- \rightarrow W^+W^-$; (b) the complete set of contributions to $e^+e^- \rightarrow W^+W^-$ in the standard model.

The fact that gauge-symmetry-violating interactions of W bosons stand out at high energies allows one to search for these interactions with great sensitivity. For example, Maalampi, Schildknecht, and Schwarzer[14] have discussed consequences of such interactions which might be made visible in studies of W pair production at LEP-II. Suzuki[15] has noted that one such interaction is already tightly constrained. Consider the term

$$\delta \mathcal{L}_{\text{eff}} = ie\,\Delta\kappa\, W^+_\mu W^-_\mu F^{\mu\nu}_{(\gamma)} , \tag{8}$$

a coupling of the W to the photon which is forbidden by $SU(2)$ gauge invariance but allowed by the gauge-invariance of electromagnetism. This term essentially provides an anomalous magnetic moment for the W boson. The interaction (8) is doubly dangerous, because it breaks not only W gauge invariance but also custodial $SU(2)$. (Direct couplings to electromagnetism, of course, do not respect custodial $SU(2)$.) The danger becomes concrete when one computes the contribution to the W mass shown in Fig. 2. The result is

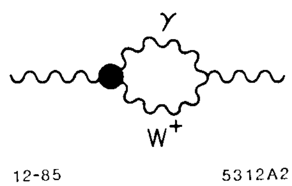


Fig. 2 A contribution to the W mass from the interaction term (8).

$$\Delta m_W^2 \sim \alpha \cdot (\Delta\kappa) \cdot \frac{\Lambda^4}{m_W^2} , \tag{9}$$

where Λ is the cutoff used in the loop integration and a standard massive vector propagator was used in the loop. There is no comparable correction to the mass of the Z^0. Thus, one finds

$$\frac{\Delta\rho}{\rho} = \left(\frac{\Delta m_W^2}{m_W^2} - \frac{\Delta m_Z^2}{m_Z^2} \right) = \frac{\alpha}{8\pi}\,\Delta\kappa\,\frac{\Lambda^4}{m_W^4} . \tag{10}$$

Prof. DiLella[16] has reported to this Symposium a new low-energy constraint on ρ:

$$\rho|_{\text{LE}} = 1.02 \pm .02 \tag{11}$$

and values of the W and Z masses which, combined with the low-energy determination of $\sin^2\theta_w$, lead to

$$\rho = \frac{m_W^2}{m_Z^2 \cos^2\theta_w|_{\text{LE}}} = 1.01 \pm .03 \tag{12}$$

If we take (10) at face value, then, this would constrain (8) to

$$\Delta\kappa < 7 \times 10^{-3} \quad (\text{for } \Lambda = 1 \text{ TeV}) . \tag{13}$$

Proponents of theories of composite W bosons must, then, take care that the interaction (8) does not appear, or appears only with a very small coefficient.

This is a very stringent requirement. For example, Kugo, Uehara, and Yanagida[17] have given a very clever formulation of the theory of composite W bosons in which this theory has a genuine, though hidden, gauge symmetry. However, the interaction (8) is allowed by this symmetry. Other approaches based on vector dominance, such as those of Fritzsch, Schildknecht, and Kögerler,[18] must also take the smallness of $\Delta\kappa$, or some equivalent statement, as an assumption.

• Composite Higgs Bosons? (Technicolor)

Let us now turn our attention from the weak vector bosons themselves to the particles which give them mass—the Higgs bosons. In contrast to the vector bosons, for which the hypothesis of composite structure is highly constrained, the Higgs bosons almost ask to be dynamically generated by a deeper theory. Aside from the gauge coupling constants (which emerge from a unique constant in grand unified theories), all of the arbitrary parameters of the standard model involve the couplings of Higgs bosons to other particles or to one another. A particularly annoying adjustable parameter is the value of the bare Higgs boson mass. It is this parameter that sets the value of the weak scale. Not only can this parameter not be determined in the standard model, but no symmetry (except perhaps supersymmetry) can keep it naturally small relative to the very deep and fundamental mass scales of Nature. It is an attractive hypothesis, then, that the weak mass scale is produced dynamically as the mass or inverse size of composite Higgs bosons.

At the first level, this program looks quite natural and achieves some easy successes. When examined more carefully, it runs into some serious, though not insurmountable, problems. I will quickly survey the present situation. For a more thorough discussion of the underlying theory, the reader should consult the reviews of Farhi and Susskind,[19] Kaul,[20] and Lane.[21]

The simplest theory of composite Higgs bosons is the original construct of Weinberg[22] and Susskind.[23] These authors postulated a new strong-interaction gauge theory, called technicolor, acting at a much larger mass scale than the conventional strong interactions. This theory contained two flavors, (U, D) of massless technifermions. They assumed that the physics of these new strong interactions was exactly that of the familiar strong interactions, scaled to the new characteristic mass. Following this analogy, the technicolor theory has an $SU(2)_L \times SU(2)_R$ chiral symmetry which should be spontaneously broken to its vector subgroup, isospin $SU(2)$. If the technifermions are coupled to the weak interactions in the conventional way, so that (U_L, D_L) form a weak doublet, the spontaneous breaking of chiral symmetry breaks the weak-interaction symmetry and gives mass to the W and Z bosons. The unbroken isospin symmetry acts as a custodial $SU(2)$ to preserve the relation $\rho = 1$. The measured masses of the W and Z are reproduced if the pion decay constant of the technicolor strong interactions takes the value

$$F_\pi = 240 \text{ GeV} ; \tag{14}$$

this corresponds to a techni–ρ meson mass of 1.8 TeV. This mechanism can be seen to work in exactly the same way, with a value of F_π half that of (14), if the technifermion doublet is made a technifamily $(\mathcal{U}, \mathcal{D}, \mathcal{N}, \mathcal{E})$, with the standard-model quantum numbers of (u, d, ν, e).

It is not a trivial problem to allow this dynamically generated Higgs couple to the familiar quarks and leptons and give them mass. To introduce this coupling, Dimopoulos and Susskind[24] and Eichten and Lane[25] postulated new gauge interactions, called extended technicolor (ETC). These new bosons would allow the constituent mass of the technifermions, generated through chiral symmetry breaking, to feed down to a mass for the ordinary fermions, by the mechanism shown in Fig. 3(a). (The technifermion mass term plays the role of the Higgs boson vacuum expectation value.) The resulting mass will be small (compared to (14)) if the ETC boson involved is heavy. Thus, the ETC bosons can be considered to contribute the the effective Lagrangian the 4-fermion vertex shown in Fig. 3(b). A four-fermion interaction has a coefficient with dimensions $(mass)^{-2}$; thus, dimensional analysis gives the following estimate of the induced masses of quarks and leptons:

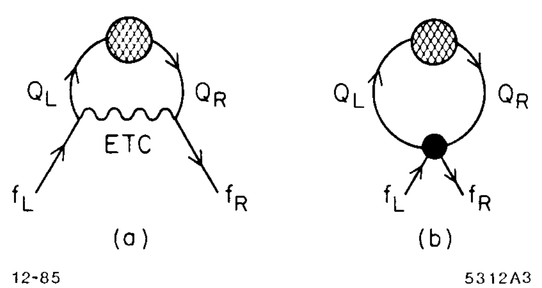


Fig. 3. The extended-technicolor mechanism for quark and lepton mass generation: (a) shows the basic diagram; (b) shows its representation in an effective Lagrangian analysis.

$$m_f = \frac{\Lambda_{TC}^3}{M_{ETC}^2} , \tag{15}$$

where Λ is the technicolor binding scale. A quark or lepton mass of 1 GeV requires a corresponding ETC boson mass of 10 TeV.

Once the theory has been fleshed out this level, however, it meets some quite nontrivial phenomenological constraints. The first of these is that theories with a technifamily predict, under quite general assumptions, the existence of light charged scalar mesons. The Higgs mechanism of vector boson mass generation requires that a massless scalar particle be absorbed by each vector boson which becomes massive, to provide its longitudinal component. In the Weinberg-Susskind model, the technicolor analogue of the pions are massless in the absence of weak interactions (because U and D are assumed massless) and become the longitudinal components of $W^{\pm}$ and Z^0 when the model is coupled to weak interactions. In technifamily models, there are effectively eight flavors and many more possible pseudoscalar mesons. The meson with content $(\bar{\mathcal{D}}\mathcal{U} + \bar{\mathcal{E}}\mathcal{N})$ becomes a component of the W^+, but the orthogonal combination $(\frac{1}{3}\bar{\mathcal{D}}\mathcal{U} - \bar{\mathcal{E}}\mathcal{N})$ remains an independent, physical boson P^+, which is massless to all orders in technicolor interactions and to the leading order in electroweak interactions. The leading contributions to the mass of this particle have been computed as follows:

$$m_{P^+}^2 = \begin{cases} (5-8\ \mathrm{GeV})^2 \\ (8-14\ \mathrm{GeV})^2 \end{cases} \qquad (\mathrm{2nd-order\ electroweak})^{25,26} \\ + \ 0\ (!) \qquad (\mathrm{1st-order\ ETC})^{27} \quad . \tag{16}$$

The two numbers given for the electroweak contribution refer to two different classes of models. Prof. Komamiya[28] has presented to this Symposium the results of searches for these particles; they are now excluded, in a clean and model-independent way* over the entire mass range allowed by (16).

I should remark parenthetically that the mass computations summarized in (16) apply only in technicolor models; in a more general context, the question of the existence of charged Higgs bosons is still open. If $m_{H^+} < m_t$, the decay $t \to b + H^+$ would be the dominant decay mode of the t quark. Thus, if the UA1 signal for t is confirmed, $m_{H^+} > 40$ GeV. If not, perhaps the t quark has a mass of 30 GeV but decays mainly to charged Higgs bosons. In this case, TRISTAN will be a Higgs factory.

The other phenomenological constraints on technicolor theories depend sensitively on the structure of the ETC couplings. If the ETC gauge group contains transitions between ordinary and techni-fermions, it must also (by its group property) include transitions linking ordinary fermions with one another. It might also include additional transitions, beyond those in technicolor itself, between technifermions. These latter transitions are relatively unconstrained; however, Appelquist, Bowick, Cohler, and Hauser[29] have noted that, if they

* unless $\mathrm{BR}(P^+ \to c\bar{b}) > 98\%$

violate custodial $SU(2)$, they can produce corrections to the ρ parameter (5) as large as 1%.

New four-fermion interactions among the familiar quarks and leptons, on the other hand, are readily observed, a situation which leads both to opportunities and to danger for this theory. The reason is that the ETC sector is precisely the part of the theory responsible for producing flavor-dependence of quark and lepton masses, and for producing flavor mixing. Thus, one must expect the induced effective 4-quark interaction to be off-diagonal in flavor. The same situation could well appear in 4-lepton couplings. Thus, we should expect that these new 4-fermion interactions mediate flavor-changing rare processes such as $K_L \to e\mu$, and $\mu \to e\gamma$, as well as $K^0 - \bar{K}^0$ and $D^0 - \bar{D}^0$ mixing, at some level. But the mass of the ETC bosons is constrained by eq. (15), and thus we know that this level is observably high[25,30,31]. The branching ratios for $K_L \to e\mu$ and $\mu \to e\gamma$, for example, are both predicted to be roughly 10^{-10}, a value which could be observed or cleanly excluded by experiments now underway at Brookhaven and Los Alamos.

Unfortunately, the simplest estimates also predict $K^0 - \bar{K}^0$ mixing amplitudes larger than the measured amplitude by a factor of 10^3. Thus, it is very important to find a realization of ETC in which this particular process is suppressed by some analogue of the GIM mechanism.[32] To show you that this is difficult but not impossible, I will review a number of proposals for achieving this suppression which have appeared in the literature.

A simple proposal is to assume that the dynamics of (d, s, b) mass generation is flavor-conserving, and that all flavor-mixing arises among (u, c, t). This idea has been put forward, in various realizations, by many authors.[33–36] This mechanism does remove $K^0 - \bar{K}^0$ mixing (along with rare K and μ decays), but it forces $D^0 - \bar{D}^0$ mixing to remain substantial. Hadeed and Holdom[36] have worked though a detailed model and predicted, in the context of this model, $\Delta m_D / m_D \sim 1 \times 10^{-13}$. This is in marginal disagreement with a new bound on $D^0 - \bar{D}^0$ mixing reported to this Symposium by the BCDMS collaboration:[37] $y < 1.2\%$, or $\Delta m_D / m_D < 0.6 \times 10^{-13}$.

A second proposal, due to Holdom,[38] assumes that the technicolor interactions are not asymptotically free. Then one should expect the particular 4-fermion interaction shown in Fig 3(b) to scale with some anomalous dimension, altering the dependence of (15) on the ETC boson mass to $(M_{ETC})^{2-\gamma}$, where γ is unknown. If $\gamma > 0$, the value of M_{ETC} required by (15) is raised, and thus ETC effects on ordinary processes are suppressed. All of the pieces of this scenario are plausible if non-asymptotically free strong-interaction gauge theories actually exist. However, the theoretical evidence for the existence of such theories is rather limited.[39]

Two more recent proposals are much closer to the spirit of GIM. Dimopoulos, Georgi, and Raby[40] have put forward a model with three ETC gauge groups $SU(N)_L \times SU(N)_U \times SU(N)_D$, coupling directly to left-handed quarks, and right-handed up and down quarks, respectively. The U and D groups can be broken to products of $U(1)$ groups at a relatively early stage, providing a mechanism for introducing large mass splittings between generations. Flavor-changing neutral-current processes, on the other hand, require the complete breaking of $SU(N)_L$ and the mixing of all three groups, and thus receives some considerable suppression. Chao and Lane[41] have proposed a scheme in which no flavor-changing neutral currents appear at tree level: The result follows from some stringent group-theoretic requirements—the various quark species q_L, u_R, d_R must belong to different but quasi-equivalent representations of ETC (*e.g.*, N and $\bar{N}$), and all bosons mediating 4-quark interactions must receive the same mass at tree level—by using the ETC couplings to break the degeneracy among possible vacuum states of the theory and then noting a special simplicity of the vacuum state chosen by this procedure. Both models have the problem that they make it difficult to generate a large mass hierarchy $m_u << m_c << m_t$. However, they signal a movement in a promising direction.

What, then, is the status of theories of technicolor? Models with a complicated technicolor sector seem to be ruled out by the nonexistence of the P^+ boson. On the other hand, the simplest Weinberg-Susskind scheme is still alive. The problem of flavor-changing neutral processes is severe but not insuperable, though the known solutions to this problem are balky and complex. To this summary, let me add one further complaint: Though technicolor seems to offer the possibility that the quark and lepton masses could be computed (in agreement with experiment, or not) from an underlying gauge theory, this promise has not yet been realized in any model. What the theory now needs, more than anything else, is a model which would allow such calculations to be done.

Before leaving the subject of technicolor, I should note a variant of it which has been advanced recently by Kaplan and Georgi[42] and explicated by Banks.[43] These authors have suggested that new 1 TeV strong interactions need not break the weak interaction symmetry directly. Consider, for example, a theory in which the group which chiral symmetry breaking leaves unbroken contains $O(4)$, which is isomorphic to $SU(2) \times SU(2)$. The unbroken group could then neatly contain $SU(2) \times U(1)$; an $O(4)$ vector of pseudoscalar mesons $(\Pi^0, \Pi^1, \Pi^2, \Pi^3)$ transforms under $SU(2) \times U(1)$ as a conventional scalar Higgs doublet. If one couples this model not only to electroweak interactions but also to an extra axial $U(1)$ boson, the pion (mass)^2 shift generated by this boson turns out to be negative and can be thought of as the negative Higgs boson (mass)^2. Then (Π^1, Π^2, Π^3) are absorbed by $W^\pm$ and Z^0 in the standard way, making these bosons massive. Π^0 becomes a neutral Higgs boson with a calculable mass, which turns out to be 150–200 GeV in specific realizations of this scenario.

• Composite Quarks and Leptons?

Finally, let us take up the question of whether the quarks and leptons might be composite states. At a first level, the motivations for imagining these particles to be composite are even more compelling than those for Higgs bosons. There are known to be many quarks and leptons; indeed, the multiplicity of these particles and their repetition of quantum numbers is one of the central mysteries of fundamental physics. Their mass spectrum is not at all understood. In fact, the failure of technicolor models to confront this problem drives us to make more far-reaching dynamical assumptions. However, the idea of composite structure within quarks and leptons quickly meets a serious and troubling dynamical question. The power of the assumption of quark and lepton compositeness depends very much on how, and how confidently, one answers this question.

To explain this basic difficulty, let me define a parameter Λ to represent the mass scale of the binding of quark and lepton constituents (preons). Equivalently, Λ^{-1} gives the physical size of the bound state. The problem is then the following: Quarks and leptons are observed to behave as pointlike particles in reactions involving momentum transfers as high as 40 GeV (at PETRA and the SPS collider). Thus, $\Lambda > 40$ GeV. On the other hand, all quarks and leptons except the t have masses much less than 40 GeV, and some have masses less than 10^{-4} times this value. Apparently, the observed quarks and leptons do not receive mass from their internal structure; more formally, the effective Lagrangian derived from the preon-binding interactions does not contain terms such as

$$\delta \mathcal{L}_{\text{eff}} = \Lambda \bar{q} q \tag{17}$$

which would give quarks and leptons masses of order Λ. We must ask, then, how such a term might be suppressed or forbidden.

To sharpen this question, let us imagine the idealized limit of a theory of composite quarks and leptons in which the composite states are exactly massless. I will discuss later how small perturbations of the dynamics can make these masses nonzero. Two general principles are known which can insure this masslessness. The first is chiral symmetry. In relativistic theory, a fermion mass term is a helicity flip operator, a mixing of left- and right-handed components of the fermion. If these left- and right-handed states have different quantum numbers under an unbroken symmetry (a chiral symmetry), that symmetry will prohibit mass generation. This scenario for the formation of massless composite fermions was put forward by 't Hooft[44] and Dimopoulos, Raby, and Susskind.[45] The assumption of an unbroken chiral symmetry comes naturally from the study of the standard model; there the weak-interaction $SU(2)$ is a chiral symmetry which prohibits any familiar fermion from acquiring mass until the Higgs field acquires its vacuum expectation value and breaks this symmetry. However, when we think of a theory with composite quarks and leptons, we must consider that this theory has strong interactions. The familiar strong interactions do not leave unbroken chiral symmetries; indeed, the strong interactions drive dynamical quark mass generation, which spontaneously breaks all such symmetries. Is is possible that a strong interactions would not cause spontaneous mass generation? 't Hooft[44] proposed a necessary condition for evading chiral symmetry breaking, the *anomaly matching condition.* Though not a sufficient condition, it is quite a stringent one. Its physical basis was reviewed, for example, in ref. 1. In the intervening time, several authors have produced extensive catalogues of the solutions to 't Hooft's condition,[46–48] defining, then, strongly-interacting gauge theories which could potentially be turned into preon models.

The second principle which can insure the masslessness of composite fermions is supersymmetry, a symmetry which connects fermions and bosons. Supersymmetry can make its beneficial influence felt in several different ways. The first mechanism proposed, by Bardeen and Visnjic,[49] postulated that supersymmetry should be spontaneously broken; each broken supersymmetry charge leads to a massless Goldstone fermion. Unfortunately, theories with many supersymmetry charges are required to contain high-spin particles; a theory with sufficient supersymmetry to produce two generations of quarks and leptons as Goldstone fermions requires spin 8. This mechanism, however, suggested another, formulated by Buchmuller, Love, Peccei, and Yanagida,[50] which is quite attractive. Imagine that, in a supersymmetric theory, an ordinary continuous symmetry is spontaneously broken. This requires the presence of a massless Nambu-Goldstone boson. The boson must have a supersymmetry partner, also massless, the *quasi-Nambu-Goldstone fermion.* By breaking a collection of continuous symmetries, one can form a multiplet of such massless fermions, which could possibly be made into quarks and leptons. This mechanism is not incompatible with the protection of fermion masslessness by chiral symmetries; indeed, as we shall see, these two mechanisms can potentially work together in a powerful way. A third mechanism, mysterious and exceptionally beautiful, appears in the superstring theory. The theory of strings provides a picture of composite quarks, leptons, and gauge bosons of striking elegance which is actually inconsistent unless it contains 10-dimensional supersymmetry, with light fermions as the supersymmetry partners of gauge bosons. I will not discuss this mechanism further, however, because it is treated in detail in Prof. Green's contribution to this Symposium.[51]

Having now reviewed the principles which might lead to the presence of light composite fermions, let us consider the basic experimental probes for such composite structure. To begin, let us return to the most general effective Lagrangian describing physics below the preon-binding scale:

$$\begin{aligned}\mathcal{L}_{\text{eff}} &= \bar{q} i \not{D} q + \bar{\ell} i \not{D} \ell + m_q \bar{q}_L \phi q_R + \ldots \\ &\qquad + \frac{\gamma_{qq}}{\Lambda^2} \bar{q}\gamma^\mu q \bar{q}\gamma_\mu q + \frac{\gamma_{q\ell}}{\Lambda^2} \bar{q}\gamma^\mu q \bar{\ell}\gamma_\mu \ell + \ldots .\end{aligned} \tag{18}$$

The kinetic terms are prescribed by $SU(3) \times SU(2) \times U(1)$ gauge invariance, as we have discussed above. The mass terms will, in general, violate symmetries of the strong-interaction theory; again, this is the explanation for the small size of these masses relative to Λ. The first terms in $\mathcal{L}_{\text{eff}}$ which are specific to the details of the underlying preon theory are the effective 4-fermion (and other dimension 6) vertices which are not present in the standard model. (A complete catalogue of possible dimension 6 structures has been compiled by Burges and Schnitzer.[52])

Experimental bounds on processes mediated by these dimension 6 operators give constraints on the size of Λ. These constraints are model-dependent to the extent that the values of the dimensionless parameters γ_{ff} can vary from model to model. The most stringent constraints will come from operators which can mediate rare flavor-changing processes; however, these operators may be forbidden by appear if the preon-binding theory conserves the flavor charges involved. It is useful to distinguish three distinct classes of preon theories, which are subject to very different constraints on Λ. The first class of theories are those which allow arbitrary flavor mixing, and, in particular, the term

$$\delta\mathcal{L}_{\text{eff}} = \frac{\gamma}{\Lambda^2} \bar{s}\gamma^\mu d \bar{s}\gamma_\mu d . \tag{19}$$

Since (19) can mediate $K^0 - \bar{K}^0$ mixing, which is very strongly constrained, these theories require $\Lambda > 10^3$ TeV. A second class of theories is that in which (19) is removed by a GIM cancellation, but the operator

$$\delta \mathcal{L}_{\text{eff}} = \frac{\gamma}{\Lambda^2} \bar{s}\gamma^\mu d \bar{e}\gamma_\mu \mu \tag{20}$$

is allowed by the symmetries of the model. Bars[47,53–55] has emphasized that this operator is difficult to remove in models in which the repetition of quark and lepton generations is natural, since (20) conserves generation number in a simple way. (20) mediates the rare decays $K \to \mu e$, $\pi \mu e$; the current upper limits on these decay rates lead to a rough limit $\Lambda > 30$ TeV.

Finally, it is possible that all dangerous flavor-changing operators are forbidden by symmetries. In this case, the preon-binding interactions should still produce effective 4-fermion couplings which are flavor-diagonal, for example,

$$\delta \mathcal{L}_{\text{eff}} = \frac{\gamma}{2\Lambda^2} \bar{e}\gamma^\mu e \bar{e}\gamma_\mu e \ . \tag{21}$$

These contact interactions will not mediate processes forbidden by symmetries of the standard model, but they will change the rates of processes for which the standard model makes quantitative predictions. Eichten, *et. al.*,[56] pointed out that the deviations produced by operators such as (21) should be relatively large, first, because they can arise from interference terms between the 4-fermion couplings and the standard model amplitudes and, second, because the 4-fermion couplings arise from strong interactions, while the standard-model contributions are suppressed by powers of α or α_s. To account properly for this latter point, they suggested the parametrization

$$\delta \mathcal{L}_{\text{eff}} = \frac{\eta_{LL} g^2}{2\Lambda^2} \bar{e}_L \gamma^\mu e_L \bar{e}_L \gamma_\mu e_L \ , \tag{22}$$

where g is taken to be of order g_ρ—$g^2/4\pi = 1$—and $\eta = \pm 1$. Over the past few years, many experiments have reported lower bounds on the value of Λ arising from this parametrization. Prof. Komamiya[28] has summarized the lower limits on Λ from PEP and PETRA experiments on Bhabha scattering, the best of which are 1–2 TeV, depending on the detailed Lorentz structure assumed for (22). PLUTO and TASSO have reported stronger limits on Λ in $e^+e^- \to \mu^+\mu^-$, which are of interest in models where e and μ share common constituents. Additional constraints on e-μ interactions have been presented by the Berkeley-TRIUMF μ decay experiment, which has reported $\Lambda > 2.9$ TeV for certain Lorentz structures.[57] UA2 has reported $\Lambda > 370$ GeV for quark-quark scattering (based on one particular, arbitrarily chosen, form for $\delta\mathcal{L}_{\text{eff}}$).[16] The general impression that one obtains from this collection of limits is that, at least for the lightest generation, the preon binding scale is well above 1 TeV.

There is a second way in which the composite structure of quarks and leptons can manifest itself experimentally. If quarks and leptons are composite, one would expect to see excited states e^* and q^* which could decay back to e or q, plus a photon or gluon. The phenomenology of these states has been discussed extensively in the literature; see refs. 58, 59 for early analyses, and refs. 60, 61 for two recent, comprehensive treatments. I will confine myself here to some brief remarks.

Like any charged fermions, e^* and q^* can be pair-produced from a γ, Z^0, $W^\pm$, or g. In principle, an f^* can also be singly produced in associated with an $\bar{f}$, as shown in Fig. 4(b). This coupling cannot have the form of a standard gauge interaction—since it is off-diagonal, it would violate gauge invariance—and so must involve extra powers of the momentum of the gauge boson, being, for example, of the form of a magnetic moment term. A magnetic moment operator has dimension 5, but it also involves a helicity flip. If the preon model contains a chiral symmetry which is only weakly broken, this interaction term must be proportional to the symmetry-breaking parameter and thus will be suppressed by roughly the same factor as the f^* mass. We would then expect Fig. 4(b) to be reasonably described by

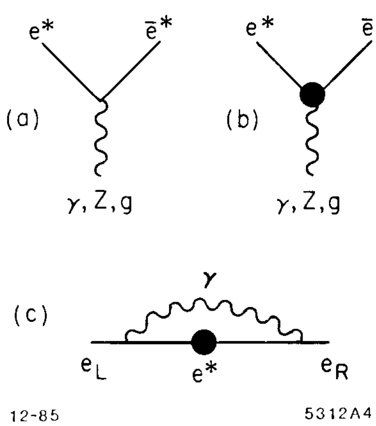


Fig. 4. Diagrams involving excited states of quarks and leptons: (a) pair production; (b) associated production; (c) a possible electron mass renormalization.

$$\delta\mathcal{L}_{\text{eff}} = e\left(\frac{m_f^*}{\Lambda^2}\right)\bar{f}^*\sigma^{\mu\nu}F_{\mu\nu}f\ . \tag{23}$$

Because this term carries a Λ^{-2}, however, it is very small if $\Lambda \sim 1$ TeV. This is, in fact, an advantage: The graph shown in Fig. 4(c) gives a renormalization of the electron mass of order

$$\alpha m_e^*(\lambda)^2\Lambda^2\ , \tag{24}$$

where, in the parametrization (23), $\lambda = (m_e^*/\Lambda)^2$. Such a small value of λ gives sufficient suppression; however, most phenomenological analyses take λ to be of order 1, replacing the quantity in parentheses in (23)by (λ/m_f^*). (There are ways to keep λ large and still suppress this diagram, such as taking (23) to involve only e_L and e_R^*.[62]) Most searches for single production of f^*'s are sensitive only to moderate values of λ. As an example, Fig. 5 shows the limits on m_e^* and λ obtained by the CELLO experiment,[28] plotted together with the estimate (23) for λ. Because of this, I give the most weight to bounds on excited quarks and leptons which are based only on pair-production. Prof. Komamiya[28] has reported new lower mass limits from PETRA on e^*, μ^*, and τ^* pair production, which, for obvious reasons, are all near 23 GeV. Ellis, Matsuda, and McKellar[63] have noted that constraints on the nuclear parity-violating potential leads to a relatively strong constraint

$$\frac{m_q^*}{\Lambda^2} = \frac{1}{3\ \mathrm{TeV}}\ ; \qquad (25)$$

it should be noted, however, that the corresponding effect in supersymmetry models[64] turns out to be naturally GIM suppressed, and that could happen here as well.

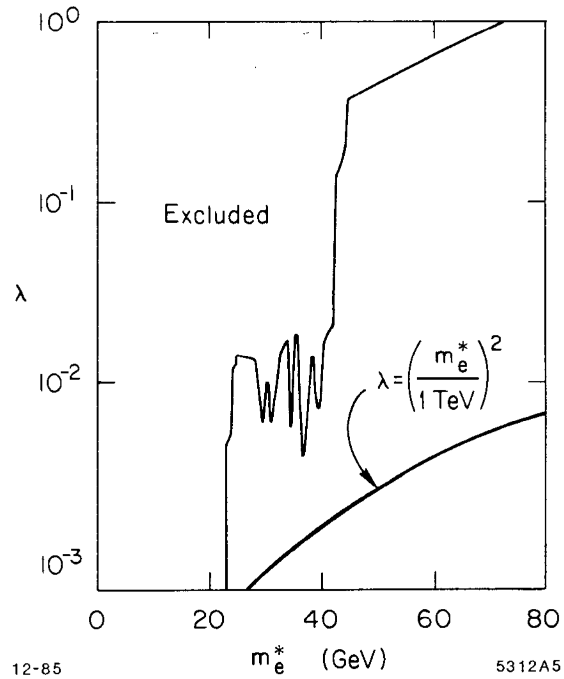


Fig. 5. Experimental limits on the mass and coupling of the m_e^*, presented by the CELLO collaboration, plotted together with the expectation for λ given by eq. (23).

If the preons inside leptons carry color, it is natural to expect that some of the excited states of leptons will be color octet particles.[65–68] The color octet e_8^* will be especially noticeable at HERA, since it can be formed as a resonance in e-g scattering. The ν_8^* has a missing-energy decay $\nu_8^* \rightarrow \nu + g$, which has made it already a target of searches at the CERN collider; UA1[69] has reported a pair production bound $m(\nu_8^*) > 60$ GeV. Models in which vector bosons are composite predict, in a similar way, color triplet vector bosons; the phenomenology of these particles has been discussed in some detail by Bauer and Streng.[70] These bosons mediate leptoquark-exchange currents which might also be visible as corrections to standard-model deep inelastic scattering. Heusch and Zerwas[71] have estimated that the NMC muon experiment at CERN will be able to search for effective interactions of this type up to $\Lambda \sim 2$ TeV. One should also expect the appearance of isoscalar weak bosons and their associated currents; Kuroda, Schildknecht, and Schwarzer have recently surveyed possible tests for such currents.[72]

I should now turn from purely phenomenological issues to a more theoretical one. The great promise of the idea of quark and lepton substructure is that having an explicit, mechanical model of fermion constituents should allow one to compute the fermion masses. I would like to explain, then, how far we have come toward realizing this promise. As a preface to this discussion, I must make two remarks. First, in any theory which encompasses the standard model, fermion mass generation will be forbidden until $SU(2) \times U(1)$ is broken. Thus, composite fermions must be born massless. They acquire their masses through their couplings to the Higgs bosons; thus, it is the magnitudes of these couplings which must we must endeavor to calculate. Second, the masses of the known quarks and leptons (even excluding neutrinos) span 5 orders of magnitude. The minimum we should ask from a scheme for fermion mass generation is that it include hierarchies in which the masses of some species are suppressed by extra powers of a small parameter.

To begin, we must ask what classes of theories can produce massless composite fermions. I have surveyed heuristic answers to this question already in ref. 1. The main progress on this question over the past few years has been negative: A number of authors have derived powerful restrictions on the appearance of massless fermions in gauge theories with strong interactions. For gauge theories of fermions and gauge bosons only, Weingarten[73] and Witten[74] have proved rigorously, under the assumption that the gauge couplings are vectorial, that the π meson, whose masslessness signals chiral symmetry breaking, is necessarily lighter than the lightest fermion. Vafa and Witten[75] have proved, under the same strong assumption, that vector (as opposed to axial-vector) flavor symmetries cannot be spontaneously broken. Thus, it is quite likely that gauge theories with vector couplings leave all fermions massive, break chiral symmetry, and leave the various flavors of fermion degenerate. The extension of these results to gauge theories with handed couplings is not at all trivial, though, and it is not unreasonable to expect that some of these theories do exhibit massless composite states. Banks and Kaplunovsky[76] managed to concoct an example of a lattice gauge theory which could be shown to have massless composite fermions for sufficiently strong gauge coupling, so this may serve to indicate at least the existence of this phenomenon.

The rigorous analysis I have just quoted does not extend straightforwardly to supersymmetric gauge theories. However, more explicit analysis of these theories has given some strong clues to their behavior,[77–79] especially in the case of vector couplings. A particularly compelling picture of this class of theories has been assembled by Affleck, Dine, and Seiberg.[80] Their argument relies on the fact that, in supersymmetric gauge theories with vector couplings, there is always some line of possible vacuum expectation values of the squark (matter boson) fields which costs no vacuum energy. Then higher-order effects on the vacuum energy along this line have the form shown in Fig. 6: To any finite order in perturbation theory, there is no effect, by virtue of a supersymmetry nonrenormalization theorem. When nonperturbative effects (for example, instantons) are taken into account, one finds a potential which is generally nonzero but which must vanish as the asymptotically free coupling vanishes for large squark vacuum expectation values. Adding to the theory a small bare squark mass m_0 produces a minimum of the potential at a large value of the squark expectation value given by an inverse power of m_0. This behavior seems almost pathological, though we will see in a moment that it can be used to advantage when this theory is embedded in a larger theory with handed couplings.

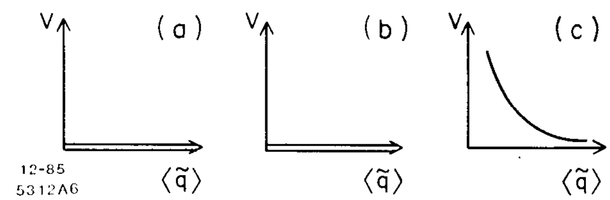


Fig. 6. The effective potential in supersymmetric gauge theories, along the line of states which at zeroth order form degenerate minima, in successive approximations: (a) zeroth order, (b) nth order perturbation theory, (c) including instantons.

Let us now assume that we can evade these constraints and form fermions as strongly-coupled bound states. How can we generate a hierarchical spectrum of fermion masses? The basic question is, what is the expansion parameter which governs this hierarchy? I will review various proposals which have appeared in the literature. Many authors have proposed that the hierarchy arises from successive powers of a small gauge coupling constant.[81–84] These authors have proposed that the chiral symmetries which insure the masslessness of the composite fermions are explicitly broken by the coupling to the preons of some weakly-interacting gauge boson. Then diagrams involving one gauge boson can change a chiral charge by 1 unit, diagrams with two gauge bosons can change this charge by 2 units, and so on. The chiral charge difference of the left- and right-handed components of a given fermion then determines the magnitude of the mass terms which is generated for this particle. Matumoto and Yamawaki[85] have shown that the expansion in the parameter $1/N_{sc}$, where N_{sc} is the number of colors in the preon-binding gauge theory, can have a similar structure and induce a mass hierarchy in the same way.

An alternative suggestion[84] has been to apply the ETC mechanism of technicolor theories, associating the spectrum of fermion masses with an m^{-2} spectrum of gauge bosons which connect the fermion constituents to the Higgs constituents. Pati[86,87] has advocated a particular realization of this idea in which there are two compositeness scales, with $\Lambda_M >> \Lambda_H$. Λ_M sets the scale of the electron and muon size, and Λ_H sets the scale of the τ and Higgs boson size. The τ can then couple directly to Higgs bosons through the Λ_H forces. The e, however, couples to the constituents of Higgs bosons only through an effective interaction generated at the scale Λ_M, as indicated in Fig. 7. This suppresses the mass of the electron (and the other fermions of the first two generations) by a factor $(\Lambda_H/\Lambda_M)^2$. Two aspects of the more detailed realization of this scheme are noteworthy for phenomenology: First, Λ_H sets the scale of the Higgs boson masses; thus we require $\Lambda_H \sim 1$ TeV. Four-e contact interactions (22) should be quite small, since these are generated at Λ_M, but $e^+e^- \to \tau^+\tau^-$ may receive large corrections. Secondly, the restrictions discussed by Bars for avoiding $K_L \to \mu e$ can be maintained in the effective theory at Λ_H but not in the full dynamics evident at Λ_M. In fact, since the e and μ generations are born at this scale, it is easy to produce an effective interaction of the form (20), with $\Lambda = \Lambda_M$. In this model, then, "BR$(K_L \to \bar{\mu}e) \geq 10^{-10}$ [is] a must".[88]

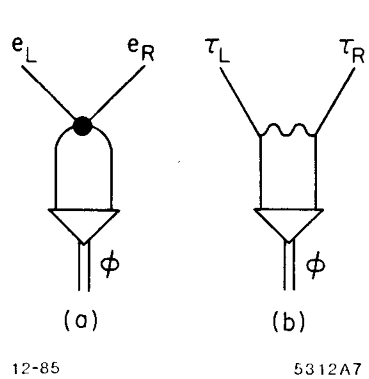


Fig. 7. Pati's mechanism for forming a hierarchy of fermion masses:
(a) e-Higgs coupling;
(b) τ-Higgs coupling.

Even more interesting hierarchical mass patterns can be obtained by applying these mechanism in supersymmetric models. If the preon-binding interactions do not break supersymmetry, some fermions will be kept massless because they are quasi-Nambu-Goldstone fermions; some of these may also be kept massless by unbroken chiral symmetries. Thus, we find double (or multiple) protection of masslessness.[89] If the various protecting symmetries are broken *sequentially*, one generates a mass hierarchy. Some alternative realizations of this idea have been studied in refs. 90–93. A recent paper by Masiero, Pettorino, Roncadelli, and Veneziano[94] brings this idea to a quite sophisticated level in constructing an almost-realistic 1-generation model with *no* ad hoc dynamical assumptions. Their scheme begins with supersymmetric Yang-Mills theory, realized in the manner suggested by Affleck, Dine, and Seiberg. They add soft supersymmetry breaking, in the form of a s-preon mass. This gives mass to the scalar partners of composite fermions, while maintaining the fermions' protection. Then they add electroweak interactions, with a gaugino mass. The gauge particle exchanges feed this mass down to the composite fermions, producing a mass spectrum which is calculable precisely in terms of (unfortunately, unknown) parameters of the composite state wave functions.

All of these schemes for generating the quark and lepton mass spectrum are somewhat crude, and none is fully predictive. But if the the quark and lepton masses have an origin in physics, we need some set of wheels and gears whose interlocking produces their observed values. The models I have just discussed give some idea, I hope, about what some basic pieces of this mechanism might be.

Digression: Precision Weak Interactions

I have reserved one important test of composite structure for a separate discussion, because it provides an example of an alternative way to probe experimentally for physics beyond the standard model.

Let me begin by discussing this specific experiment, the measurement of the muon $(g-2)$. The rapport between the experimental value of the muon $(g-2)$ and that computed in the standard model is now expressed by[95] :

$$a_\mu - a_\mu|_{\text{st. model}} = 38 \pm 85 \pm 20 \ \times 10^{-10} \ ; \tag{26}$$

the first error is that of the experimental determination, the second that of the calculated standard-model value. To see what effect compositeness of the muon might have on this agreement, let me parametrize the anomalous magnetic moment due to the muon's internal structure by writing an effective interaction analogous to (23)[96,97]

$$\delta \mathcal{L}_{\text{eff}} = e\left(\frac{m_\mu}{\Lambda^2}\right)\bar{\mu}\sigma^{\mu\nu}F_{\mu\nu}\mu \ . \tag{27}$$

(26) implies that the parameter Λ in this equation is bounded by

$$\Lambda < 720 \text{ GeV} \qquad (90\% \text{ confidence}) \ . \tag{28}$$

This bound is roughly as strong as the bounds on Λ arising from searches for 4-fermion contact interactions. (One should keep in mind that the various parameters Λ which I have defined may differ by factors of 2 or π, in a model-dependent way.)

It is interesting to ask whether this bound can be improved in the near future. The value of a_μ in the standard model,

$$a_\mu|_{\mathrm{st.\ model}} = 11\ 659\ 202\ (20) \ \times 10^{-10} , \tag{29}$$

is already known to 1.7 parts per million; this is the error recorded in (26). Of this, the pure QED contribution is extremely well understood, contributing an error of only 0.26 ppm. The contribution of weak interactions is readily calculated; the whole effect turns out to be 1.7 ppm. The dominant source of error, accounting for 1.6 ppm, is the set of diagrams which involve hadronic contributions to the vacuum polarization and hadronic light-by-light scattering. The latter set of diagrams have recently been reexamined and computed cleanly by Kinoshita, Nizic, and Okamoto.[98] The contributions from vacuum polarization, though, require experimental knowledge of the cross section for $e^+e^- \rightarrow hadrons$ at low energies. New experiments have allowed a substantial decrease in the uncertainty in this term,[98] and there is room for further improvement. The direct determination of a_μ remains at the value given by the CERN $(g-2)$ experiments,[99]

$$a_\mu = 11\ 659\ 240\ (85) \ \times 10^{-10} , \tag{30}$$

leaving an error of 7.3 ppm. However, a group led by Hughes has made a proposal for a new experiment[95] that it claims could decrease this error to 0.3 ppm. Such an experiment could actually measure the weak-interaction contribution to the muon $(g-2)$, and could also probe for muon substructure at the scale of several TeV.

This analysis contains a lesson which is more generally applicable. In this example, but also more generally in weak-interaction physics, contributions from new physics beyond the standard model are of the same order of magnitude as W boson loop corrections. Thus, precision tests of weak-interaction radiative corrections become, as well, probes for such new physics. In general, experiments sensitive enough to make these tests can see effects of new particles with mass up to several hundred GeV or values of Λ above 1 TeV.

What other experiments are likely to probe so sensitively? To discuss this question, I should remind you that experiments which involve the weak interactions directly are very sensitive to the three basic parameters of the weak interactions—the two coupling constants g and g' and the Higgs field vacuum expectation value $\langle\phi\rangle$. Before trying to test the theory, one needs to determine these parameters precisely; this requires measurement of three physical quantities. Two of these are provided by α and G_μ; the third will be in place when m_Z is measured to tens of MeV at SLC and LEP. One way to express the accuracy of this measurement is that it will determine $\sin^2\theta_w$ to a precision $\delta\sin^2\theta_w \sim 10^{-4}$. In making such a statement, however, one must realize that $\sin^2\theta_w$ is a derived quantity which is defined by a convention. An appropriate convention is set out and discussed in ref. 100.

Once the three basic quantities α, G_μ, and m_Z have been measured, further precision measurements can provide nontrivial tests of the standard model or probes beyond it. The two most promising quantities for further precision experiments are the W boson mass m_W[101] and the polarization asymmetry in Z^0 production A_{LR}.[100] Experiments now being planned should accurately measure both of these quantities within the next eight years: The D∅ experiment at the Tevatron has been designed to determine m_W with minimum systematic error; the D∅ proposal claims that $\delta m_W \sim 50$ MeV may be achieved.[102] The quantity A_{LR}, which may be defined as

$$A_{LR} = \left.\frac{\sigma(e_L^- e^+ \to \mu\bar{\mu}) - \sigma(e_R^- e^+ \to \mu\bar{\mu})}{\sigma(e_L^- e^+ \to \mu\bar{\mu}) + \sigma(e_L^- e^+ \to \mu\bar{\mu})}\right|_{s=m_Z^2}, \tag{31}$$

can be measured directly at the SLC. With an improved source of polarized electrons, one might reach $\delta A_{LR} \sim 3\times 10^{-3}$. Both of these determinations would correspond to $\delta\sin^2\theta_w \sim$ a few $\times 10^{-4}$. A better way to understand the sensitivity of these measurements, however, is to ask whether they have reached the level at which small deviations from the standard model become apparent. To show that this is so, I display in Fig. 8 the effect on these two quantities of two such modifications. Even a single heavy quark doublet added to the standard model can make its presence felt.

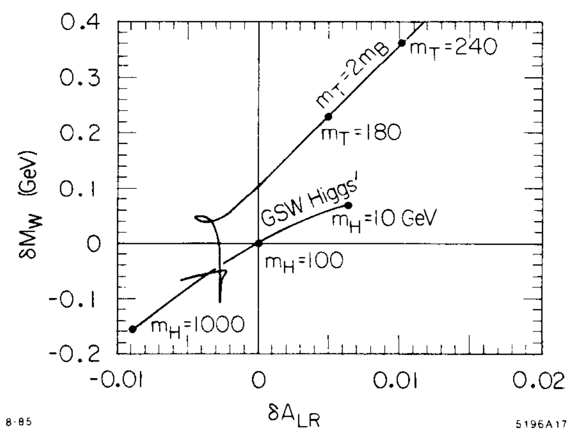


Fig. 8. Effect on the quantities m_W and A_{LR} of two modifications of the standard model: variation of the Higgs boson mass up to 1 TeV, and addition of a heavy quark doublet with $m_T = 2m_B$.

The connection between the various weak-interaction quantities I have discussed is, of course, not without theoretical uncertainty. At the level of these experiments, however, the only important theoretical uncertainty comes, again, in our knowledge of the hadronic part of the QED vacuum polarization.[103] For these weak-interaction probes, as for the muon $(g-2)$, we need improvements in the experiments which determine this quantity, the measurement of the cross-section for e^+e^- annihilation into any hadrons, especially in the region $1 < \sqrt{s} < 2$ GeV.

New Strong Interactions at 1 TeV

The last part of this lecture will deal with a topic which is new to this series of conferences, but which should be of increasing importance in the future. It has already been recognized as a crucial issue in the plans in the United States and Europe for very high energy proton colliders. This is the possibility of new strong interactions at the 1 TeV scale, and the question of how these interactions manifest themselves experimentally. Most of the models I described in the first part of this lecture require such new physics. What I wish to point out here is that, almost independently of the details of the specific model in question, new 1 TeV strong interactions are likely to be visible, and even striking, to experiments of high enough energy to reach their natural scale.

• Minimal Model

I will begin by discussing the simplest model of TeV strong interactions, the standard model with a large value of the Higgs boson mass. There are two ways to understand why such a model should be strongly interacting. The first is to recall that the renormalizable Higgs field potential energy $V = -m^2\phi^*\phi + \frac{1}{2}\lambda(\phi^*\phi)^2$ yields a Higgs field vacuum expectation value

$$\langle\phi\rangle = \left(\frac{m^2}{\lambda}\right)^{\frac{1}{2}} . \tag{32}$$

The physical Higgs boson mass, equal to $\sqrt{2}m$, can then be much larger than $\langle\phi\rangle = 240$ GeV only if the ϕ coupling constant λ is much larger than 1. Alternatively, one may recall the logic of Lee, Quigg, and Thacker[13] which I discussed at the start of this lecture: the process $W^+W^- \to W^+W^-$ requires gauge-theory cancellations to respect unitarity in perturbation theory. In particular, it requires a contribution from diagrams with Higgs scalar exchange. But this contribution may be selectively suppressed by making the Higgs scalar heavy.

Unitarity is then violated in perturbation theory if the Higgs is sufficiently heavy, $m_H > (8\pi\sqrt{2}/3G_F) = 1$ TeV. In that case, higher orders of perturbation theory must be of the same size as the leading-order terms.

It is striking that strong interactions may appear so naturally in the Higgs sector. Once they appear there, it is straightforward for them to couple into other sectors of the theory, especially the sector of the weak vector bosons, to which the Higgs fields give mass. The influence of Higgs boson dynamics on high-energy interactions of W bosons is summarized elegantly in the *equivalence theorem* shown in Fig. 9(a)[104,105,13]. Since a massless vector boson is purely transverse, while a massive vector boson has also a longitudinal component, a W boson can become massive only by absorbing one additional degree of freedom, a charged scalar component of the Higgs field. This field then loses its separate identity. If one quantizes weak-interaction theory in the Unitarity gauge, it formally disappears from the theory entirely, though in the Feynman–'t Hooft gauge it is kept and summed together with other unphysical (ghost and timelike-polarized) modes which give cancelling contributions. The equivalence theorem states that a remnant of this field remains visible in physical amplitudes, that the cross-section for producing a longitudinally-polarized W boson at an energy high relative to the W mass is equal to the cross section for producing the original, ephemeral scalar. An especially easy case to understand is the one shown in Fig. 9(b). If Higgs bosons have strong interactions, then the graph on the right-hand side of this relation is a strong-interaction amplitude, unsuppressed by powers of α. In Feynman-'t Hooft gauge, this graph is a contribution to the cross-section for longitudinal W production, and it is the only such contribution which contains no explicit α. A somewhat less trivial example of the theorem is the result that the sum of the graphs of Fig. 1(b), evaluated for $s >> m_W^2$, is just equal to the amplitude for $e^+e^- \to \phi^+\phi^-$, via a virtual photon and Z^0. Recently Chanowitz and Gaillard[106] have given an elegant new proof of this theorem

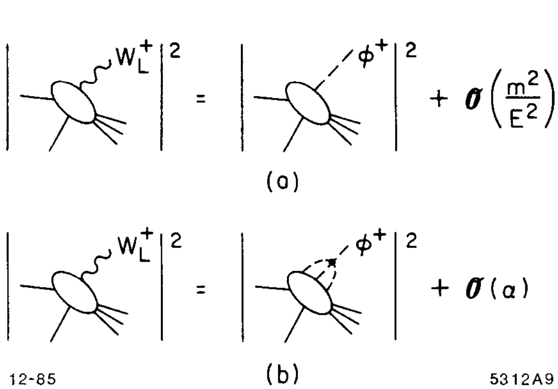


Fig. 9. The equivalence theorem relating the amplitudes for emitting longitudinally-polarized W bosons and Higgs scalars: (a) shows the general result; (b) gives a special case which is very easy to understand.

which is sophisticated enough to apply to weak and electromagnetic production processes and to situations where several W_L's are produced.

The equivalence theorem tells us that when the Higgs sector becomes strongly interacting, the W and Z bosons do as well. This has two consequences of special importance, which I would now like to explore. The first of these is a new process for producing Higgs bosons (and W and Z pairs) in high-energy collisions. This process was first discussed by Jones and Petkov;[107] its importance has been emphasized more recently by Kane, Repko, and Rolnick[108] and, especially, by Cahn and Dawson.[109] (Some more detailed studies of this process are given in refs. 110–112.) The mechanism is that shown in Fig. 10(a): colliding quarks or leptons bremsstrahlung W or Z bosons, which then fuse to form a Higgs resonance. If the massive bosons are longitudinally polarized, the factor of α coming from their coupling to fermions is compensated by the strong-interaction coupling to the Higgs. This process dominates the more conventional process of gluon-gluon fusion (shown in Fig. 10(b)), which has the disadvantageous dependence $(m_t/m_H)^2$ for large Higgs boson masses, for m_H larger than a few hundred GeV. A comparison of the two processes, in terms of total cross-section for Higgs production in pp collisions, is shown in Fig. 11. The magnitude of Higgs boson production in high-energy pp collisions is shown in another way in Fig. 12, taken from the comprehensive study of Eichten, Hinchliffe, Lane, and Quigg.[113] This figure shows the magnitude of the process $pp \to H \to W^+W^-$. Since the Higgs boson is strongly coupled to W and Z bosons, these also form the dominant decay products of the Higgs almost as soon as the Higgs mass moves above W^+W^- threshold. For $m_H >> m_W$,

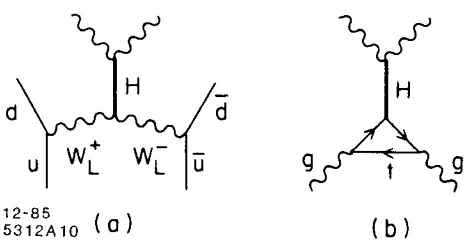


Fig. 10. Processes for Higgs boson production in very high energy collisions: (a) vector boson fusion, (b) gluon fusion.

$$\Gamma_H \simeq \frac{3\sqrt{2}}{32\pi} G_F m_H^3 . \tag{33}$$

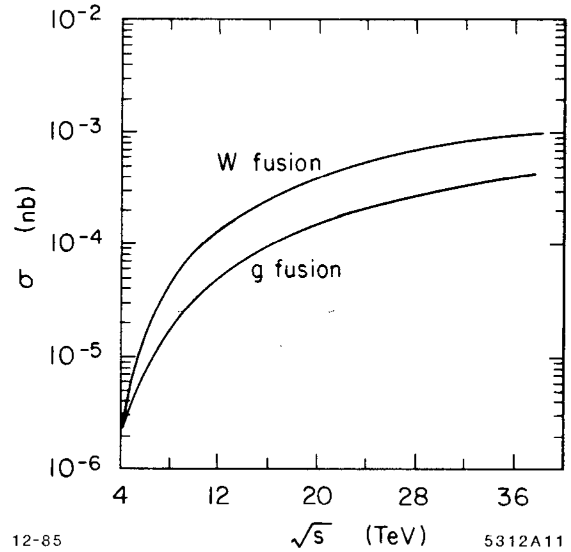


Fig. 11. Total cross-sections for Higgs boson production in very high energy pp collisions, comparing the W fusion and gluon fusion mechanisms, for $m_H = 7m_W$, from ref. 109.

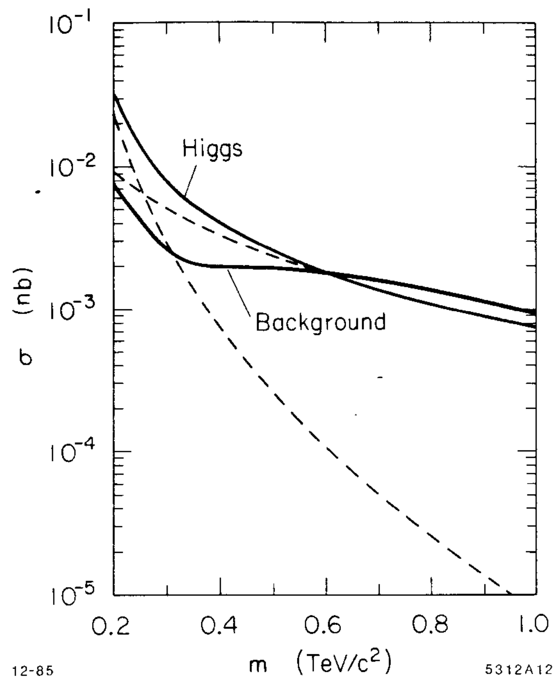


Fig. 12. Cross-section for Higgs boson production, and subsequent decay to W^+W^-, in 40 TeV pp collisions, as a function of the Higgs mass, computed for $m_t = 30$ GeV and using a rapidity cut $|y| < 2.5$. The dashed curves show the two components of the production process. The background shown is the physics background in its narrowest sense, the cross-section for the process $q\bar{q} \to W^+W^-$, integrated over W pair masses within $\Gamma_H/2$ of m_H. (from ref. 113)

Over a large range of masses, this is a great advantage, because it gives the Higgs boson a spectacular signature of a W or Z pair summing to a fixed invariant mass. This promising situation eventually disappears, however, because the width (33) increases so rapidly with m_H. For $m_H \sim 1$ TeV, $\Gamma_H \sim 500$ GeV, and the Higgs peak disappears into the W^+W^- continuum.

For Higgs masses above 1 TeV, it seems fruitless to search for the Higgs boson as a sharp resonance. But if the Higgs boson is this heavy, its physical effects might be even more interesting. This is the second consequence of the equivalence theorem, that if the Higgs sector is strongly interacting, these interactions will produce distortions in the cross-sections for longitudinal W and Z production. I would like to discuss two different effects of this nature. The first is the effect on the simple annihilation process e^+e^- (or $q\bar{q}$) $\to W^+W^-$ of W-W final-state interactions. By helicity conservation, the annihilation of light fermions requires $J = 1$ and so also, by Bose statistics of the W_L's, $I = 1$, where the isospin I is custodial $SU(2)$. The Higgs sector interactions are expected to be strongest in the $I = J = 0$ channel, but that channel is not available here. Nevertheless, the process is highly constrained kinematically, especially in e^+e^- annihilation, where one has the further advantage that the W production amplitude is precisely calculable. Fig. 13 gives some idea of the visibility of corrections to the lowest-order amplitude. It shows that the cross-section for

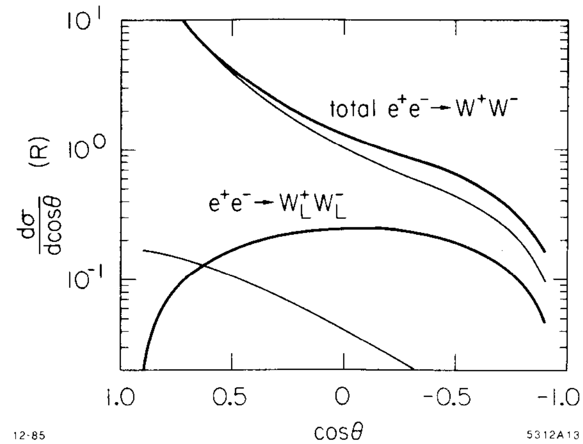


Fig. 13. Differential cross-sections (in units of R) for $e^+e^- \rightarrow W^+W^-$ into states of definite helicity, for $\sqrt{s} = 1$ TeV. The various curves show the total production, and the partial cross-sections for $W_T^+W_T^-$, $W_L^+W_L^-$, and $W_L^{\pm}W_T^{\mp}$.

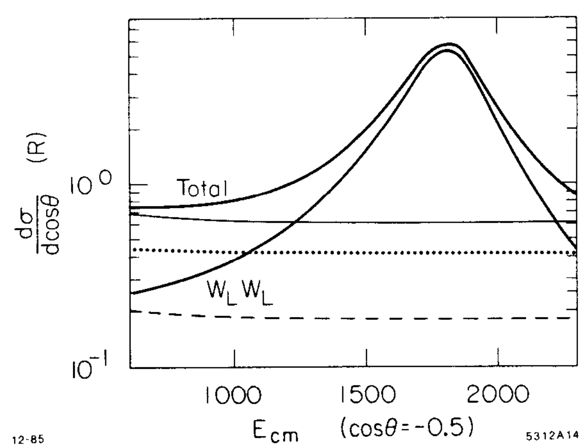


Fig. 14. Helicity-state cross-sections (in units of R) for $e^+e^- \rightarrow W^+W^-$ at $\cos\theta = -0.5$ in the simplest technicolor model, showing the effect of the technicolor ρ meson resonance.

$e^+e^- \rightarrow W_L^+W_L^-$ is roughly $\frac{1}{4}$ of the total W pair production in the backward hemisphere. If one can concentrate on backward production, even small corrections to the $W_L^+W_L^-$ production amplitude can become visible. The simplest assumption, that one should extrapolate upward and unitarize the low-energy form of the W-W (or ϕ-ϕ) interaction, leads to a disappointingly small effect, only a few percent correction to the total backward rate.[114] On the other hand, if the strong Higgs interactions are those of the Weinberg-Susskind technicolor model,[22,23] the rho meson of the technicolor interactions appears in this channel and one finds the enormous effect shown in Fig. 14. A similar, but hardly so dramatic, effect should be visible in very high energy pp collisions.[113] A second mechanism of W and Z pair production is the continuum analogue of the process of Fig. 10(a), the scattering of bosons radiated from fermion lines. (We have noted that Fig. 10(a) is actually not distinguishable from this continuum for $m_H >$ 1 TeV.) This process has the disadvantage that the center-of-mass energy of the boson pair is not defined by the kinematics. However, it has the advantage that it involves longitudinal W and Z bosons dominantly, and that it can access the $I = J = 0$ channel. Chanowitz and Gaillard[106] have made a detailed study of strong-interaction effects in this process; a sample of their results is shown in Fig. 15. Similar effects should appear in TeV-energy e^+e^- annihilation;

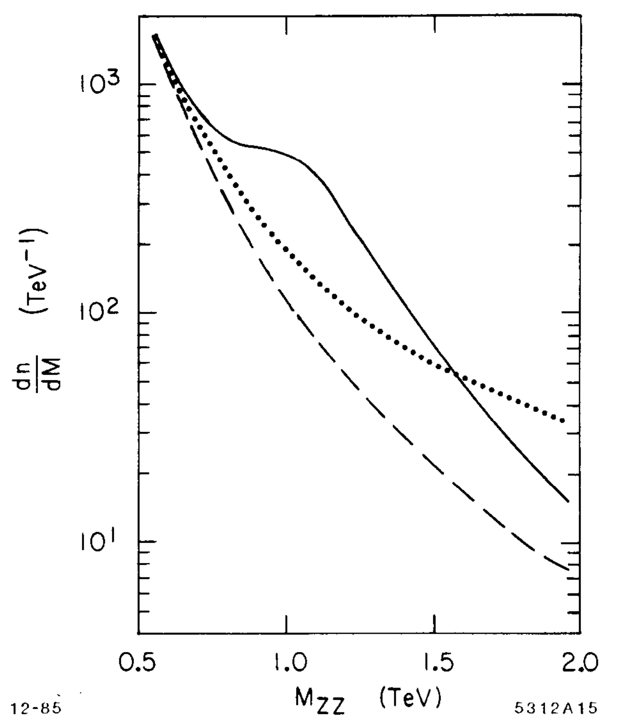


Fig. 15. Cross-sections for Z^0 pair production in 40 TeV pp collisions, as a function of the Z^0 pair invariant mass M. The dashed curve is the result of $q\bar{q}$ annihilation. The dotted and solid curves include Z radiation, with the Z_L-Z_L interaction given, respectively, by extrapolation of the low-energy formula and by inclusion of a Higgs scalar resonance at a mass of 1 TeV. (from ref. 106)

a large background identified by Dawson and Rosner[115] can probably be eliminated by considering (as is done in ref. 106) only boson pairs with mass above 1 TeV.[116] This question, however, needs further study.

I must conclude this discussion of strong interaction effects in weak boson pair production, however, with a word of warning. It is not at all obvious that W bosons can be distinguished from gluon jets in TeV-energy collisions. If the W is to be identified in its hadronic decay mode, assuming that its mass is measured to ± 10 GeV, the signal sits three orders of magnitude below the background from QCD jets.[117] If one selects events with one leptonic W decay, the background is still comparable in size to the signal, even with well-chosen cuts.[118–120] Leptonic decays of the Z^0 should tag these particles quite specifically, but requiring this decay mode cuts deeply into the rate of processes which will already be rare. It is possible that this background problem might be a fundamental limitation our ability to do weak boson physics with pp colliders. I hope, though, that one can simply find a more effective trick for plucking out W's and Z's. I commend this very important problem to your attention.

- **Embellishments**

The simple example of the standard model with a heavy Higgs boson already shows some interesting new physics accessible to TeV-energy colliders. However, one should consider even more seriously the possibility that TeV-energy physics will be governed by one of the more detailed scenarios described in the first part of this lecture. In general, adding more structure at 1 TeV gives more remarkable phenomena to be observed. In this section, I will give a few examples of phenomena which appear in these more detailed schemes, drawing both on technicolor and on composite-fermion models.

I have already noted that the techni-ρ can appear as a dramatic resonance in W and Z pair production. Let me consider, as well, two effects more peculiar to technicolor models. The first is the possible presence of long-lived techni-baryons. Rubakov and collaborators[121] have pointed out that the techni-baryon may be stable with respect to perturbative interactions but unstable via a weak-interaction barrier penetration effect. This effect requires that the technibaryon decays predominantly to 12-fermion final states

$$B \rightarrow 3\begin{pmatrix} u \\ d \end{pmatrix} + 3\begin{pmatrix} c \\ s \end{pmatrix} + 3\begin{pmatrix} t \\ b \end{pmatrix} + \begin{pmatrix} \nu_e \\ e \end{pmatrix} + \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix} + \begin{pmatrix} \nu_\tau \\ \tau \end{pmatrix} , \tag{34}$$

producing a dramatic signature. Unfortunately, a Skyrme model estimate predicts that the decay rate is large enough for (34) to be observed only if $m_B > 12$ TeV. The second effect is the production of ETC bosons. The relation (15) implies that, if the t quark mass is 40 GeV, the ETC boson which produces this quark mass is very light, of order 1 TeV. These ETC bosons can then be pair-produced at high-energy colliders. Arnold and Wendt[122] have studied the production of these bosons in 40 TeV pp collisions and found some quite unexpected features. Since ETC bosons have technicolor interactions, they form techni-hadronic bound states. Gluon-gluon collisions dominantly produce the lowest few bound states. These then decay weakly to t quarks and technihadrons. A dominant decay chain is shown in Fig. 16; this leads to the process

$$g + g \rightarrow (t + (\bar{t} + Z^0)) , \tag{35}$$

where each set of parentheses denotes a combination of definite invariant mass. The cross section for the chain of processes turns out to be large—1 nb for $m_{\rm ETC} = 1$ TeV, decreasing to 10 pb for $m_{\rm ETC} = 1.5$ TeV.

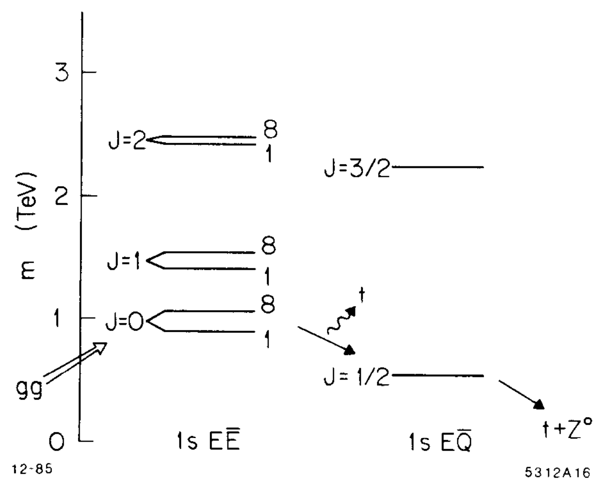


Fig. 16. Mechanics of ETC boson production in pp collisions. The figure shows the spectrum of ETC-pair and ETC-techniquark bound states (from ref. 122) for $m_{\rm ETC} = 1$ TeV, and the observable processes which link them.

Theories with composite quarks and leptons lead to even more dramatic phenomena, since there the electrons and quarks which are the primary probes become strongly interacting. If the contact interactions such as (22) are barely visible at $\sqrt{s} = 100$ GeV, they become dominant already at energies of several hundred GeV. This leads to rising cross-sections

$$\sigma(ff \to ff) \sim \frac{s}{\Lambda^4} \tag{36}$$

which develop, for $\sqrt{s} \sim \Lambda$ into a sequence of resonances. In e^+e^- annihilation, one should see the whole evolution in total cross-sections, as is illustrated in Fig. 17. In pp collisions, one probes a large range of $q\bar{q}$ center-of-mass

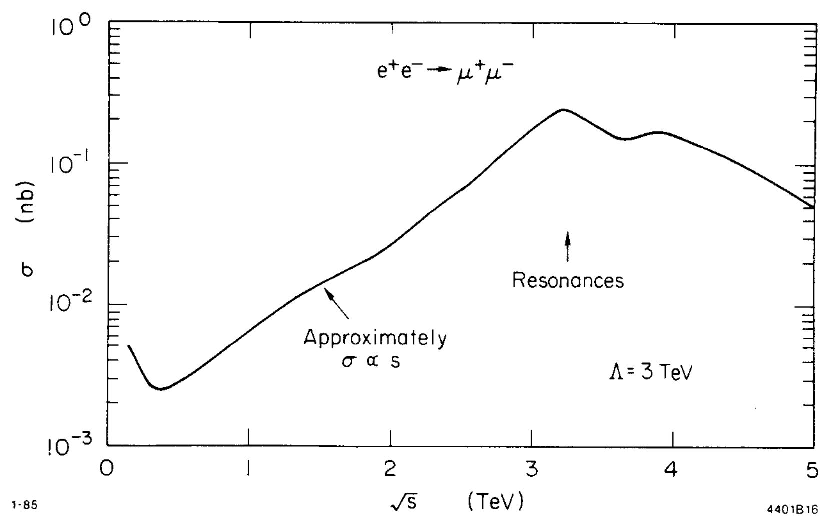
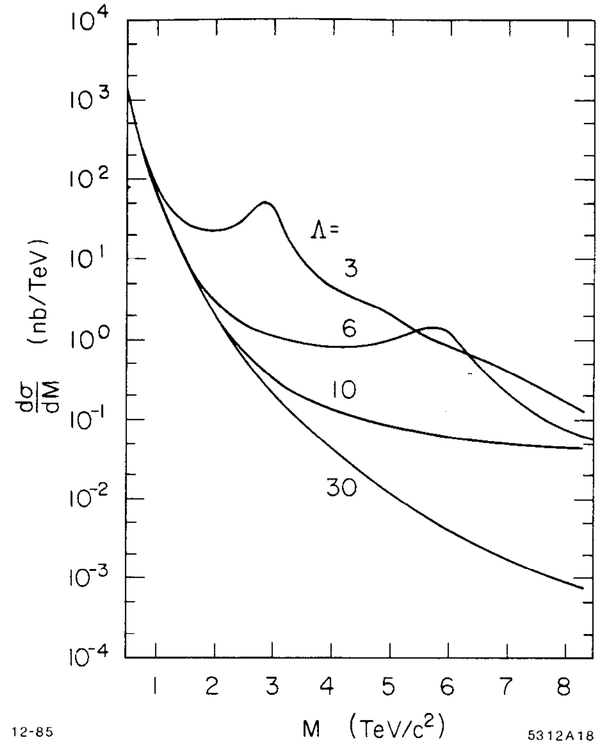


Fig. 17 Energy-dependence of the cross-section for $e^+e^- \to \mu^+\mu^-$ in a model of composite leptons with $\Lambda \sim 3$ TeV.

energies simultaneously, but this jumble of processes can be disentangled by looking in the jet or Drell-Yan $p_\perp$ spectrum, or in the spectrum of two-jet invariant masses. Detailed numerical studies of this phenomenon[113,123,124] show that one can easily pick out the rise of the cross-section for large-momentum transfer processes. For some choices of the parameters, such as that shown in Fig. 18, the resonances of the preon-binding interactions are also visible.

Fig. 18. Effect of 4-quark contact interactions on the 2-jet invariant mass spectrum in 40 TeV pp collisions, for one set of parameters chosen in ref. 123.



Conclusions

By now, we have come quite far into the realm of speculation. But it is important to remember that we were led to this point in a logical way, arguing from the basic question of explaining the masses of quarks, leptons, and gauge bosons which is the central unsolved problem posed by the success of the standard model. What do we learn from these wild ideas?

Theorists should learn not to be embarrassed by models with many visible moving parts. All of the models I have discussed are highly mechanical, and none are particularly elegant. But I feel it is important to make even contrived calculations of the fermion mass spectrum, as long as these calculations use real physical mechanisms and come to definite conclusions. How else can we make any progress at all?

Experimenters should learn to be patient. I have noted at various points in this lecture that they should continue experiments sensitive to Λ, ℓ^*'s, and q^*'s, and that they should continue searching for rare processes, such as $K \to \mu e$, and for Higgs bosons, charged and neutral. This may be a frustrating pursuit for a long time, especially when one searches at 100 GeV energies for TeV-energy phenomena. But, eventually, it is these searches that will yield the concrete information which makes the next level of physics clear.

Finally, we should all begin to plan for the era of TeV-energy colliders, an era which may be less than a decade away. At these energies, we could well see dramatic changes from our current picture of the fundamental interactions. Perhaps we can anticipate those changes, or perhaps we will be stunned when they appear. In either case, we can hardly be complacent with our current successes while we still have so much to learn.

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DISCUSSION

D. Schildknecht, Universität Bielefeld

I have a comment on your comment on my comment [on Komamiya's lecture]. Let us just assume that the W boson is of composite nature, as in one part of your talk you did. Then the first conclusion you would draw would be that the known bosons will have brothers and sisters, like an isoscalar boson or an excited boson, which for phenomenological reasons should probably have a mass above 300 or 500 GeV. Their couplings are of course determined by the overlap function between the quarks and leptons and their excited states; these should come out to be of the order of magnitude of or smaller than the coupling of the ground state and thus should be of the order of $g^2/4\pi$ or the fine structure constant. That was the comment I made this morning, and I do not think that this is a completely unreasonable picture.

M. Peskin

Actually, I am very confused as to how you can make that coupling so small, if these particles begin as strongly interacting bosons, the analogues of the ρ in the real strong interactions. Maybe you can get a reduction by $1/N_c$ from $g_\rho^2/4\pi = 2$. But how can you get a reduction of a factor of 100?

D. Schildknecht

You see, the coupling in such a composite model should be determined by the overlap. What you are raising is perhaps a question about how to make a concrete model even for the W boson.

M. Peskin

Absolutely. In my lecture, I talked about a problem for the composite W and Z models which has nothing to do with this comment you have just raised. But this problem has also disturbed me a lot about the composite W and Z models. The coupling has to be so weak, and it's hard to understand that in a real dynamical model.

J. Pati, University of Maryland

I have a small theoretical comment and one phenomenological comment. You listed some of the possible difficulties. They are really only facets of a possible theoretical difficulty for composite models. One of these is the question of the composite gauge bosons. Then there is the Weingarten type of result, that the mass of the composite boson is less than the mass of the composite fermion.

M. Peskin

The one I emphasized is only proven in vector-like theories.

J. Pati

Yes. The other is the Vafa-Witten one, which is also particular to a vector-like theory like QCD.

It seems to me that all of these must sum up somehow so as to be avoided in a natural composite model. If you think of the compositeness scale as being very much higher—which is different from what Schildknecht might be referring to; I might place the compositeness scale for W and Z even up to the grand unification scale—then there is local supersymmetry as opposed to global supersymmetry. Then the Weinberg-Witten constraint does not apply. And, by the same token, also the Vafa-Witten and Weingarten type of results do not apply. So it appears that local supersymmetry, which is naturally motivated for reasons other than that composite models have to work, has a natural role to play here. This only means that one can reverse the No-Go theorems; more work is needed to show that the desirable result really follows. But at least it is interesting that the No-Go theorems do not apply. Would you have any comment on that?

M. Peskin

In fact, no. You are absolutely right that the No-Go theorems do not apply in this case. One ought to understand these theorems better, to see if one can prove that there are counterexamples to them there. I think that is a very important problem.

J. Pati

I am going on to make a phenomenological comment. The process $e^+e^- \rightarrow \tau^+\tau^-$ is extremely important for the class of models in which the electron family and the muon family would have small size but the τ and the 4th τ' family have large size. By the same token, it is important to do these experiments with quarks and antiquarks of the electron family going into $\tau^+\tau^-$, as in the proton-antiproton experiments. There, of course, it would be very important to have vertex detectors to select out the τ channel properly.

H. Schnitzer, Brandies University

We have shown many years ago, based on some technical features of analyticity and something called Mandelstam counting, that, in the electroweak gauge theory, gauge bosons lie on Regge trajectories.[125] If there are light Higgs bosons, there are ordinary recurrences of theses bosons which are heavy. Then, unequivocably, high-spin friends of the W will appear as J resonances of the ρ type. They should be looked for, perhaps in WW scattering, in terms of angular distributions rather than just bumps.

M. Peskin

That is quite interesting; I should make a comment on it. The calculation that I showed for e^+e^- reactions were done in a simple and straightforward way without taking this effect into account. However, when you do such calculations in proton-proton scattering, you must be very careful, because you sample very higher effective values of $\sqrt{s}$ as you move toward $x = 1$ for the partons. The parton distributions are falling but the cross-sections may be rising [due to eq. (36)], so you have to use a unitarized formula. In the calculations which I discussed, Bars and his collaborators did something very pretty. To represent quark-quark scattering, they used the Veneziano formula, with the trajectories shifted to that mass values you would expect from thinking about the physics of these composite models. As you see [in Fig. 18], the first recurrence can be quite prominent. If you look closely, you can also see, in the top curve of this figure, the second recurrence. It's a little hard to see, because it is rather wide if you put in a realistic value for the width. But, in principle, these results reflect exactly the physics you are talking about.

J. Ellis, CERN

I have two sneaky questions. First, in the supersymmetric model of Masiero, Pettorino, Roncadelli, and Veneziano that you mentioned, we have just one generation. How complicated does the theory have to be, if one want to have three generations with realistic mixing? That was the easy question.

The difficult question is: From the point of view of the composite models that you have discussed, which should be more interesting, a 40 TeV center-of-mass hadron-hadron collider, or a 2 TeV center-of-mass e^+e^- collider?

M. Peskin

To the first question, I actually don't know the answer. It is relatively easy to include generations by enlarging the fundamental gauge group which appears in this model from $SU(6)$ to, I believe, $SU(4N+2)$. But I don't think they know how to get a realistic mixing pattern, and I don't know how either.

In answer to your second question, I have shown that, to some extent, these two devices are complementary. So, from a purely theoretical point of view, it is hard to judge which one is preferable. [I assume that some solution can be found to the background problems for observing WW scattering in pp collisions.] From the practical point of view, the technology for TeV-energy e^+e^- colliders is still quite far away from a realistic design. Clearly, we should be working to change that situation.