Technicolor Evolution

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This talk describes how modern theories of dynamical electroweak symmetry breaking have evolved from the original minimal QCD-like technicolor model in response to three key challenges: R_b, flavor-changing neutral currents, and weak isospin violation.

1. Introduction

In order to understand the origin of mass, we must find both the cause of electroweak symmetry breaking, through which the W and Z bosons obtain mass, and the cause of flavor symmetry breaking, by which the quarks and leptons obtain their diverse masses and mixings. The Standard Higgs Model of particle physics, based on the gauge group $SU(3)_c \times SU(2)_W \times U(1)_Y$ accommodates both symmetry breakings by including a fundamental weak doublet of scalar ("Higgs") bosons $\phi = \binom{\phi^+}{\phi^0}$ with potential function $V(\phi) = \lambda\left(\phi^\dagger\phi - \frac{1}{2}v^2\right)^2$. However the Standard Model does not explain the dynamics responsible for the generation of mass. Furthermore, the scalar sector suffers from two serious problems. The scalar mass is unnaturally sensitive to the presence of physics at any higher scale Λ (e.g. the Planck scale), as shown in fig. 1. This is known as the gauge hierarchy problem. In addition, if the scalar must provide a good description of physics up to arbitrarily high scale (i.e., be fundamental), the scalar's self-coupling (λ) is driven to zero at finite energy scales as indicated in fig. 1. That is, the scalar field theory is free (or "trivial"). Then the scalar cannot fill its intended role: if $\lambda = 0$, the electroweak symmetry is not spontaneously broken. The scalars involved in electroweak symmetry breaking must therefore be a party to new physics at some finite energy scale – e.g., they may be composite or may be part of a larger theory with a UV fixed point. The Standard Model is merely a low-energy effective field theory, and the dynamics responsible for generating mass must lie in physics outside the Standard Model.

This talk focuses on Dynamical Electroweak Symmetry Breaking, an approach that evades the hierarchy and triviality problems because the scalar states involved in electroweak symmetry breaking are manifestly composite at scales not much above the electroweak scale $v \sim 250$ GeV. . The prototypical model of this kind is Technicolor, originally conceived as having dynamics modeled on those of QCD. To build a minimal technicolor model, one starts with the Standard Model, removes the Higgs doublet, and adds an asymptotically free (technicolor) gauge force, e.g. based on the group $SU(N)_{TC}$, and a set of massless technifermions which feel this new force. The electroweak charges of the technifermions are chosen so that the formation of a technifermion condensate will break the electroweak symmetry to its electromagnetic subgroup. The simplest choice is to include two flavors of technifermions, of which the right-handed components, U_R and D_R, are weak singlets while the left-handed members form a weak doublet $(U, D)_L$. The Lagrangian for the technifermions, like that of massless up and down quarks, possesses a global

Figure 1: (left) Naturalness problem: $M_H^2 \propto \Lambda^2$. (right) Triviality: $\beta(\lambda) = \frac{3\lambda^2}{2\pi^2} > 0$.

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Figure 2: (left) Top-technifermion scattering mediated by a heavy ETC gauge boson. (right) Technifermion condensation creates the top quark mass.

chiral symmetry $SU(2)_L \times SU(2)_R$. At a scale $\Lambda_{TC} \sim 1$ TeV, the technicolor coupling g_{TC} becomes strong, causing the technifermions to condense: $\langle \bar{U}U + \bar{D}D \rangle \neq 0$. The condensate breaks the technifermions' chiral symmetries to the vector subgroup $SU(2)_{L+R}$; the Nambu-Goldstone bosons of this symmetry breaking are called technipions Π_T, in analogy with the pions of QCD. Because of the technifermions' electroweak quantum numbers, the condensate also breaks $SU(2)_W \times U(1)_Y$ to $U(1)_{EM}$, and the technipions become the longitudinal modes of the W and Z. The logarithmic running of the strong technicolor gauge coupling renders the low value of the electroweak scale (i.e. the gauge hierarchy) natural in these theories, while the absence of fundamental scalars obviates concerns about triviality.

Even a minimal technicolor sector, as described above, should yield visible signatures in collider experiments. For example, the technihadron spectrum, like the QCD hadron spectrum, will include vector resonances like techni-rho and techni-omega states (ρ_T, ω_T). These should contribute to the rescattering of longitudinal electroweak bosons radiated from initial-state quarks at the LHC. For relatively light states $M_{\rho_T} \sim 1TeV$, a peak might be visible in the invariant mass spectrum for production of $W+Z$ where both bosons decay leptonically [1, 2]. An even sharper signal due to mixing of the ρ_T with electroweak bosons could also be present in direct $q\bar{q} \to W_L Z_L$ processes [3]. While one might also have hoped to detect technicolor through enhanced non-resonant W^+W^+ scattering, the signal is neither large nor kinematically distinct from the background [4]. At a linear $e+e^-$ collider with $\sqrt{s} = 1.5$ TeV and $\mathcal{L} = 200$ fb^{-1}, a 1 TeV vector resonance could make its presence felt in $e^+e^- \to W^+W^-\bar{\nu}\nu$ (but not in $e^+e^- \to ZZe^+e^-$), while the non-resonant "low-energy theorem" contribution would, again, be undetectable [5]. The Z-boson form factor would also be sensitive to a ρ_T with a mass up to a few TeV [6]. Finally, at a 4 TeV muon collider with $\mathcal{L} = 200$ fb^{-1}, the gauge-boson rescattering process $\mu^+\mu_- \to W^+W^-X$ (ZZX) with hadronically-decaying W's (Z's) will be sensitive to a technirho of up to 2 TeV (1 TeV); same-sign W production from like-sign muon beams shows only a featureless increase in the WW invariant mass [7].

In order to generate masses and mixings for the quarks and leptons, it is necessary to couple them to the source of electroweak symmetry breaking. The classic way of doing this is by extending the technicolor gauge group to a larger extended technicolor (ETC) group under which the ordinary fermions are also charged. When ETC breaks to its technicolor subgroup at a scale $M > \Lambda_{TC}$, the gauge bosons coupling ordinary fermions to technifermions acquire a mass of order M. At the scale Λ_{TC}, a technifermion condensate breaks the electroweak symmetry as described earlier, and the quarks and leptons acquire mass because the massive ETC gauge bosons couple them to the condensate. The top quark's mass, e.g., arises when the condensing technifermions transform the scattering diagram in fig. 2 (left) into the top self-energy diagram shown at right. Its size is

$$m_t \approx (g_{ETC}^2/M^2)\langle T\bar{T}\rangle \approx (g_{ETC}^2/M^2)(4\pi v^3) \quad . \tag{1}$$

Thus M must satisfy $M/g_{ETC} \approx 1.4$ TeV in order to produce $m_t = 175$ GeV.

While this mechanism works well in principle, it has proven difficult to construct a complete model that can accommodate the wide range of observed fermion masses while remaining consistent with precision electroweak data. Three key challenges have led model-building in new and promising directions. First, the dynamics responsible for the large value of m_t must couple to b_L because t and b are weak partners. How, then, can one obtain a predicted value of R_b that agrees with experiment? Attempts to answer this question have led to models in which the weak interactions of the top quark[8, 9, 10, 11] (and, perhaps, all third generation fermions) are non-standard. Second, while creating large fermion masses m_f requires M_{ETC} to be of order one TeV, suppressing flavor-changing neutral currents demands that M_{ETC} be several orders of magnitude higher. Attempts to resolve this conflict have led to the idea that the technicolor gauge dynamics may have a small beta-function: $\beta_{TC} \approx 0$ [12]. Such "walking technicolor" models often include light

Figure 3: Fermion currents coupling to the weak-singlet ETC boson that generates m_t.

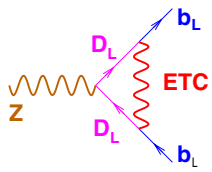

Figure 4: Direct correction to the $Zb\bar{b}$ vertex from the ETC gauge boson responsible for m_t in a commuting model.

technihadrons with distinctive signatures [13]. Third, despite the large mass splitting $m_t \gg m_b$, the value of the rho parameter is very near unity. How can dynamical models accommodate large weak isospin violation in the $t - b$ sector without producing a large shift in M_W? This issue has sparked theories in which the strong (color) interactions of the top quark[14] (and possibly other quarks[15]) are modified from the predictions of QCD.

In the remainder of this talk, we explore the ways in which modern theories of dynamical electroweak symmetry breaking have evolved in response to these issues in flavor physics.

2. The R_b Challenge

In classic extended technicolor models, the large value of m_t comes from ETC dynamics at a relatively low scale M_{ETC} of order a few TeV. At that scale, the weak symmetry is still intact so that t_L and b_L function as weak partners. Moreover, experiment tells us that $|V_{tb}| \approx 1$. As a result, the ETC dynamics responsible for generating m_t must couple with equal strength to t_L and b_L. While many properties of the top quark are only loosely constrained by experiment, the b quark has been far more closely studied. In particular, the LEP measurements of the $Zb\bar{b}$ coupling are precise enough to be sensitive to the quantum corrections arising from physics beyond the Standard Model. As we now discuss, radiative corrections to the $Zb\bar{b}$ vertex from low-scale ETC dynamics can be so large that new weak interactions for the top quark are required to make the models consistent with experiment.[8, 16, 17]

To begin, consider the usual ETC models in which the extended technicolor and weak gauge groups commute, so that the ETC gauge bosons carry no weak charge. In these models, the ETC gauge boson whose exchange gives rise to m_t couples to the fermion currents[16, 17]

$$\xi \left(\bar{\psi}_L^i \, \gamma^\mu \, T_L^{ik}\right) \; + \; \xi^{-1} \left(\bar{t}_R \, \gamma^\mu \, U_R^k\right) \tag{2}$$

where ξ is a Clebsh of order 1 (see fig. 3). Then the top quark mass arises from technifermion condensation and ETC boson exchange as in fig. 2, with the relevant technifermions being U_L and U_R.

Exchange of the same[16, 17] ETC boson causes a direct (vertex) correction to the $Z \to b\bar{b}$ decay as shown in fig. 4; note that it is D_L technifermions with $I_3 = -\frac{1}{2}$ which enter the loop. This effect reduces the magnitude of the $Zb\bar{b}$ coupling by an amount governed by the size of the top quark mass

$$\delta g_L = \frac{e}{4 \sin\theta \cos\theta} \left(\frac{g^2 v^2}{M^2}\right) \approx \frac{e}{4 \sin\theta \cos\theta} \frac{m_t}{4\pi v} \tag{3}$$

where we have used the relationship between M_{ETC} and m_t from eqn. 1.

The shift in the coupling directly affects the ratio of Z decay widths $R_b \equiv \Gamma(Z \to b\bar{b})/\Gamma(Z \to$ hadrons), such that the fractional change in R_b is proportional to δg_L. Moreover, oblique and QCD corrections to the decay widths cancel in the ratio, up to factors suppressed by small quark masses. One finds[16, 17]

$$\frac{\delta R_b}{R_b} \approx -5.1\% \cdot \xi^2 \cdot \left(\frac{m_t}{175\text{GeV}}\right) \tag{4}$$

Such a large shift in R_b is excluded[18] by the data. Then the ETC models whose dynamics produces this shift are likewise excluded.

This suggests one should consider an alternative class of ETC models[8] in which the weak group $SU(2)_W$ is embedded in G_{ETC}, so that the weak bosons carry weak charge. Embedding the weak interactions of all quarks in a low-scale ETC group would produce masses of order m_t for all up-type quarks. Instead, one can extend $SU(2)$ to a direct product group $SU(2)_h \times SU(2)_\ell$ such that the third generation fermions transform under $SU(2)_h$ and the others under $SU(2)_\ell$. Only $SU(2)_h$ is embedded in the low-scale ETC group; the masses of the light fermions will come from physics at higher scales. Breaking the two weak groups to their diagonal subgroup ensures approximate Cabibbo universality at low energies. The electroweak and technicolor gauge structure of these non-commuting models is sketched below[8]:

$$\begin{array}{rcl} G_{ETC} & \times & SU(2)_{light} \times U(1) \\ & \downarrow & f \\ G_{TC} & \times & SU(2)_{heavy} \times SU(2)_{light} \times U(1)_Y \\ & \downarrow & u \\ G_{TC} & \times & SU(2)_{weak} \times U(1)_Y \\ & \downarrow & v \\ G_{TC} & \times & U(1)_{EM} \end{array} \tag{5}$$

Due to the extended gauge structure, three sets of electroweak gauge bosons are present in the spectrum: heavy states W^H, Z^H that couple mainly to the third generation, light states W^L, Z^L resembling the standard W and Z, and a massless photon $A^\mu = \sin\theta[\sin\phi\, W^\mu_{3\ell} + \cos\phi\, W^\mu_{3h}] + \cos\theta X^\mu$ coupling to $Q = T_{3h} + T_{3\ell} + Y$. Here, ϕ describes the mixing between the two weak groups and θ is the usual weak angle. The coupling of Z_L to quarks differs from the standard model value by $\delta g_L = (c^4/x)T_{3\ell} - (c^2 s^2/x)T_{3h}$. This reduces R_b[8]

$$\frac{\delta R_b}{R_b} \approx -5.1\% \, [\sin^2\phi \frac{f^2}{u^2}] \tag{6}$$

where the term in square brackets is $\mathcal{O}(1)$.

At the same time, there is still a contribution to R_b from the dynamics that generates m_t. The ETC boson responsible for m_t now couples weak-doublet fermions to weak-singlet technifermions (and vice versa) as in fig. 5. The radiative correction to the $Zb\bar{b}$ vertex is as in fig. 4 except that the technifermions involved are now U_L with $T_3 = +\frac{1}{2}$. As a result, the shifts in δg_L and R_b have the same size as the results in eqns. (3) and (4) but the opposite sign.[8]. Overall, then, the ETC and ZZ' mixing contributions to R_b in non-commuting ETC models have equal magnitude and opposite sign, enabling R_b to be consistent with experiment. The key element that permits a large m_t and a small value of R_b to co-exist is the presence of non-standard weak interactions for the top quark[8]. This is something experiment can test, and has since been incorporated into models such as topflavor[10, 11] and top seesaw.[19, 20]

There are several ways to test whether the high-energy weak interactions have the form $SU(2)_h \times SU(2)_\ell$. One possibility is to search for the extra weak bosons. The bosons' predicted effects on precision electroweak data gives rise to the exclusion curve[9] in fig. 6. Low-energy exchange of Z^H and W^H bosons would cause apparent four-fermion contact interactions; LEP limits on $eebb$ and $ee\tau\tau$ contact terms imply[21] $M_{Z^H} \geq 400$ GeV. Direct production of Z^H and W^H at Fermilab is also feasible; a Run II search for $Z^H \to \tau\tau \to e\mu X$ will be sensitive[21] to Z^H masses up to 650 - 850 GeV. Another possibility is to measure the top quark's weak interactions in single top production. Run II should measure the ratio of single top and single charged lepton cross-sections $R_\sigma \equiv \sigma_{tb}/\sigma_{\ell\nu}$ to $\pm 8\%$ in the W^* process.[22, 23, 24] A number of systematic uncertainties, such as those from parton distribution functions, cancel in the ratio. In the Standard

Figure 5: Fermion currents coupling to the weak-doublet ETC boson that generates m_t in non-commuting ETC models.

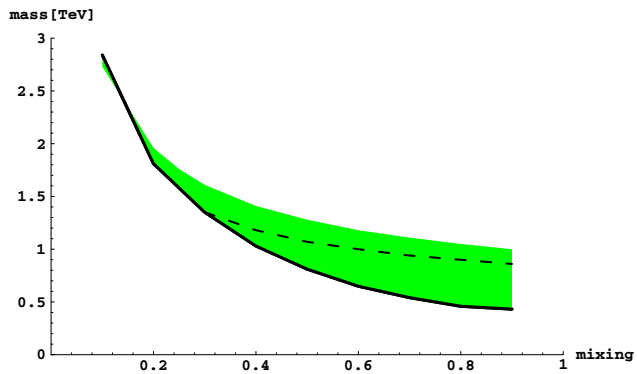

Figure 6: FNAL Run II single top production can explore the shaded region of the $M_{W'}$ vs. $\sin^2\phi$ plane.[25] The area below the solid curve is excluded by precision electroweak data.[9] In the shaded region R_σ increases by$\geq$ 16%; below the dashed curve, by $\geq$ 24%.

Model, R_σ is proportional to the square of the Wtb coupling. Non-commuting ETC models affect the ratio in two ways: mixing of the W_h and W_ℓ alters the W^L coupling to fermions, and both W^L and W^H exchange contributes to the cross-sections. Note that the ETC dynamics which generates m_t has no effect on the Wtb vertex because the relevant ETC boson does not couple to b_R. Computing the total shift in R_σ reveals (see fig. 6) that Run II will be sensitive [25] to W^H bosons up to masses of about 1.5 TeV.

3. The FCNC Challenge

In order for extended technicolor to produce the wide range of observed fermion masses, it is necessary for ETC dynamics to couple differently to like-charge fermions belonging to different generations. This causes ETC boson exchange to contribute to flavor-changing neutral currents, a potential source of severe constraints on model-building.

For example, in the neutral Kaon system, exchange of ETC bosons contributes to the square of the $K_L K_S$ mass difference by an amount of order

$$(\Delta M_K^2)_{ETC} \simeq \frac{g_{ETC}^2 \,\mathrm{Re}(\theta_{sd}^2)}{2M_{ETC}^2} \tag{7}$$

where f_K is the Kaon decay constant and M_K is the average neutral Kaon mass. Because the experimental upper bound on the Kaon mass difference is $\Delta M_K < 3.5 \times 10^{-12}$ MeV [18], one may deduce

$$\frac{M_{ETC}}{g_{ETC}\sqrt{\mathrm{Re}(\theta_{sd}^2)}} > 600\mathrm{TeV} \tag{8}$$

Remembering that the fermion mass m_f produced by exchange of an ETC boson of mass M_{ETC}

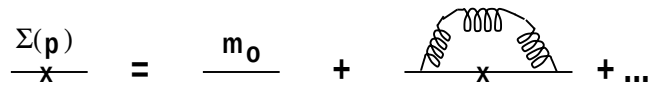

Figure 7: Gap equation for generation of dynamical technifermion mass $\Sigma(p)$.

and coupling g_{ETC} scales as $m_f \sim g_{ETC}^2/M_{ETC}^2$, we see that the limit (8) makes it difficult[1] for ETC to produce fermion masses much larger than an MeV.

In order to address this issue, let us revisit the origin of fermion masses in ETC models in a little more detail[2]. The expression for the dynamically-generated fermion mass m_f is

$$m_f \approx \frac{g_{ETC}^2}{M_{ETC}^2}\langle\overline{U}U\rangle_{ETC} \tag{9}$$

where the condensate is evaluated at the ETC scale. In previous numerical estimates of the sizes of fermion masses, we have used the approximation

$$\langle\overline{U}U\rangle_{ETC} \approx \langle\overline{U}U\rangle_{TC} \approx 4\pi F_{TC}^3. \tag{10}$$

More generally, however, the condensate scales as

$$\langle\overline{U}U\rangle_{ETC} = \langle\overline{U}U\rangle_{TC}\exp\left(\int_{\Lambda_{TC}}^{M_{ETC}}\frac{d\mu}{\mu}\gamma_m(\mu)\right) \tag{11}$$

If the technicolor gauge dynamics resemble those of QCD, the value of the anomalous dimension γ_m is small over the integration range, so that the integral is negligible and the approximation (10) holds.

This immediately raises the question: what if the technicolor coupling instead runs slowly from Λ_{TC} up to M_{ETC}? In other words, what would happen if the technicolor beta function were small, $\beta_{TC} \approx 0$, making the coupling "walk" instead of running like the QCD coupling?

The answer lies in the dynamics by which masses for ordinary fermions and technifermions arise. Consider the Schwinger-Dyson gap equation for the dynamical technifermion mass $\Sigma(p)$ in the rainbow approximation (fig. 7). The phenomenologically interesting solutions to the gap equation are those manifesting chiral symmetry breaking: those for which $\Sigma(p) \neq 0$ even if the bare mass m_o vanishes. Detailed analysis [12, 29, 30, 31] has shown that a chiral symmetry breaking occurs only when the value of $\alpha_{TC}(\Lambda_{TC})$ approaches the critical value $\alpha_c \equiv \pi/3C_2(R)$. Since the anomalous dimension $\gamma_m(\mu)$ can be written as $1-\sqrt{1-\alpha_{TC}(\mu)/\alpha_c}$, a large value of $\alpha_T C(\Lambda_{TC})$ implies that $\gamma_m(\Lambda_{TC}) \approx 1$. This gives the starting point for the integration in eqn. (11). If $\beta_{TC} \approx 0$, then $\alpha_{TC} \sim \alpha_c$ and $\gamma_m \sim 1$ persist up to the scale M_{ETC}. This enhances the integral in eqn. (11) and enables ETC to generate fermion masses as large as m_s or m_c:

$$m_{q,l} = \frac{g_{ETC}^2}{M_{ETC}^2} \times \left(\langle\overline{T}T\rangle_{ETC} \cong \langle\overline{T}T\rangle_{TC}\frac{M_{ETC}}{\Lambda_{TC}}\right) \tag{12}$$

where $M_{ETC}/\Lambda_{TC} \sim 100-1000$.

The small technicolor beta function that produces enhanced fermion masses can arise in models with many technifermions in the standard vector representation of the technicolor gauge group or in models with fermions in several different technicolor representations. In either case, the chiral symmetry-breaking sector is enlarged relative to that of minimal technicolor models. As a result, one expects a proliferation of technipions and small technipion decay constants $f_{TC} \ll v$. At first glance, it appears that the models will suffer from unacceptably large contributions to S (because of the large number of technifermions) and from the presence of many light pseudo-Nambu-Goldstone bosons (PNGBs) which have not been observed. However, the nature of the

[1] An even more stringent limit may be found in [26].
[2] This discussion was inspired by refs. [27, 28].

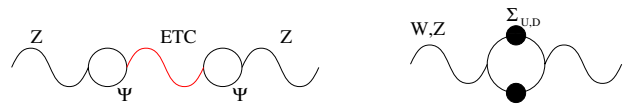

Figure 8: ETC contributions to $\Delta\rho$: (left) direct, from gauge boson mixing (right) indirect, from technifermion mass splitting.

strong walking-technicolor dynamics must be taken into account. First, QCD is no longer a reliable guide for the estimation of contributions to S; the walking models have a new pattern of resonance masses (possibly a tower of ρ_T and ω_T states), more flavors, and fermions in non-vectorial gauge representations [26]. In the absence of compelling estimates of S in walking gauge theories, S does not provide a decisive test of these models. Second, the walking dynamics which enhances the technifermion condensate also enlarges the masses of the PNGBs

$$F_{TC}^2 M_{\pi_T}^2 \approx \frac{g_{ETC}^2}{M_{ETC}^2} \left(\langle \overline{T} T \rangle_{ETC} \right)^2 \tag{13}$$

enabling them to meet current experimental constraints.

The phenomenological signatures of walking technicolor have been studied in models known as "lowscale technicolor"[32, 32, 33]. The primary signals exploit the contrasting effects of walking on the masses of the vector mesons ρ_T and ω_T (which are reduced) and those of the technipions π_T (which are enhanced). Lighter technivector mesons are more readily produced in colliders, and if the technipion masses are enhanced enough to close the $\rho_T \to \pi_T \pi_T$ and $\omega_T \to 3\pi_T$ decay channels, the technivectors will quite visibly to final states including electroweak gauge bosons. For instance [33], if one takes the number of weak-doublets of technifermions to be $N_D \approx 10$ in order to induce walking, the technivector meson masses are reduced to $M_{\rho_T} \approx M_{\omega_T} \approx 2v/\sqrt{N_D} \approx$ 200 GeV. At the same time, the effects of walking tend to raise the mass of the π_T to over 100 GeV. The dominant technirho decays will then be to one π_T plus one electroweak gauge boson or to two electroweak bosons. The π_T are expected to decay to $f\bar{f}$ pairs through ETC couplings, making decays to heavy flavors dominate.

Signatures of low-scale technicolor would be visible at both hadron and lepton colliders. Current limits have been summarized by M. Narain in these Proceedings (or see [18]); these include searches for $\rho_T \to W\pi_T$ followed by $\pi_T \to b\bar{b}, c\bar{c}b\bar{c}, c\bar{b}$, for pair-production of technipions with leptoquark quantum numbers, and for production of ρ_T and ω_T decaying to lepton pairs through mixing with electroweak gauge bosons. Future experiments will, naturally, have greater reach. For example, a 200 GeV muon collider could resolve [34] the peaks of even nearly degenerate ρ_T and ω_T in the process $\mu^+\mu^- \to \rho_T, \omega_T \to e^+e^-$. As another example, the LHC, technirhos with masses up to 500 or more GeV would provide a visible signal in 30 fb^{-1} of ATLAS data through decays to a WZ pair which then decay to leptons [35]. Summaries of the expected reach of various technicolor searches at Tevatron Run II and the LHC may be found in refs. [18] and [36], respectively.

4. The $\Delta\rho$ Challenge

At tree-level in the Standard Model, $\rho \equiv M_W^2/M_Z^2 \cos^2\theta_W \equiv 1$ due to a "custodial" global $SU(2)$ symmetry relating members of a weak isodoublet. Because the two fermions in each isodoublet have different masses and hypercharges, however, oblique radiative corrections to the W and Z propagators alter the value of ρ. The one-loop correction from the (t,b) doublet is particularly large because $m_t \gg m_b$. Experiment[18] finds $|\Delta\rho| \leq 0.4\%$, a stringent constraint on isospin-violating new physics.

Dynamical theories of mass generation like ETC must break weak isospin in order to produce the large top-bottom mass splitting. However, the new dynamics may also cause additional, large contributions to $\delta\rho$. Direct mixing between and ETC gauge boson and the Z (fig. 8) induces the dangerous effect[37, 38]

$$\Delta\rho \approx 12\% \cdot \left(\frac{\sqrt{N_D} F_{TC}}{250 \text{ GeV}} \right)^2 \cdot \left(\frac{1 \text{ TeV}}{M_{ETC}/g_{ETC}} \right)^2 \tag{14}$$

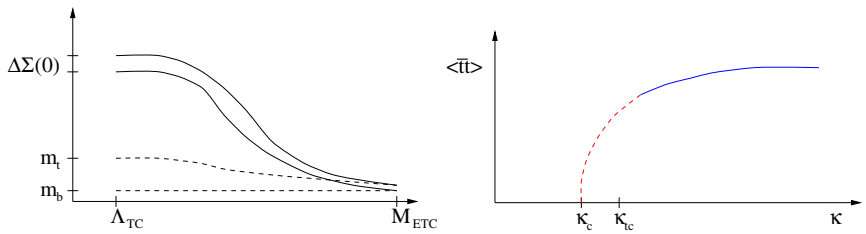

Figure 9: (left) ETC and new top dynamics generate masses for technifermions, t and b. (right) Second-order phase transition forms a top condensate for $\kappa > \kappa_c$.

in models with N_D technifermion doublets and technipion decay constant F_{TC}. To avoid this, one could make the ETC boson heavy; however the required $M_{ETC}/g_{ETC} > 5.5\text{TeV}(\sqrt{N_D}F_{TC}/250\text{ GeV})$ is too large to produce $m_t = 175$ GeV. Instead, one must obtain $N_D F_{TC}^2 \ll (250\text{GeV})^2$ by separating the ETC sectors responsible for electroweak symmetry breaking and the top mass. A second contribution comes indirectly[39] through the technifermion mass splitting: $\Delta\rho \sim (\Sigma_U(0) - \Sigma_D(0))^2/M_Z^2$, as in fig. 8. Again, a cure[14, 31, 40, 41, 42, 43, 44, 45, 46, 47, 48] is to arrange for the t and b to get only part of their mass from technicolor. As sketched in fig. 9, suppose M_{ETC} is large and ETC makes only a small contribution to the fermion and technifermion masses. At a scale between M_{ETC} and Λ_{TC} new strong dynamics felt only by (t,b) turns on and generates $m_t \gg m_b$. The technifermion mass splitting is small, $\Delta\Sigma(0) \approx m_t(M_{ETC} - m_b(M_{ETC}) \ll m_t$, and no large contributions to $\Delta\rho$ ensue.

The realization that new strongly-coupled dynamics for the (t,b) doublet could be so useful has had a dramatic effect on model-building. Models in which some (topcolor[14, 40, 41, 42, 43, 44, 45, 46, 47, 48]) or even all (top mode [49, 50, 51, 52, 53], top seesaw[54, 55]) of electroweak symmetry breaking is due to a top condensate have proliferated. One physical realization of a new interaction for the top is a spontaneously broken extended gauge group called topcolor[14, 40, 41, 42, 43, 44, 45, 46, 47, 48]: $SU(3)_h \times SU(3)_\ell \to SU(3)_{QCD}$. The (t,b) doublet transforms under $SU(3)_h$ and the light quarks, under $SU(3)_\ell$. Below the symmetry-breaking scale M, the spectrum includes massive topgluons which mediate vectorial color-octet interactions among top quarks: $-(4\pi\kappa/M^2)(\bar{t}\gamma_\mu\frac{\lambda^a}{2}t)^2$. If the coupling κ lies above a critical value ($\kappa_c = 3\pi/8$ in the NJL[56, 57] approximation), a top condensate forms (fig. 9). For a second-order phase transition, $\langle\bar{t}t\rangle/M^3 \propto (\kappa - \kappa_c)/\kappa_c$, so the top quark mass generated by this dynamics can lie well below the symmetry breaking scale; so long as M is not too large, the scale separation need not imply an unacceptable degree of fine tuning.

A more complete model incorporating these ideas is topcolor-assisted technicolor[14, 40, 41, 42, 43, 44, 45, 46, 47, 48] (TC2). The symmetry-breaking structure is:

$$\begin{aligned} G_{TC} &\times SU(3)_h \times SU(3)_\ell \times SU(2)_W \times U(1)_h \times U(1)_\ell \\ &\downarrow \quad M \geq 1 \text{ TeV} \\ G_{TC} &\times SU(3)_{QCD} \times SU(2)_W \times U(1)_Y \\ &\downarrow \quad \Lambda_{TC} \sim 1 \text{ TeV} \\ G_{TC} &\times SU(3)_{QCD} \times U(1)_{EM} \end{aligned} \tag{15}$$

Below the scale M, the heavy topgluons and Z' mediate new effective interactions[14, 15, 40, 41, 42, 43, 44, 45, 46, 47, 48, 58, 59, 60, 61, 62] for the (t,b) doublet

$$-\frac{4\pi\kappa_3}{M^2}\left[\overline{\psi}\gamma_\mu\frac{\lambda^a}{2}\psi\right]^2 - \frac{4\pi\kappa_1}{M^2}\left[\frac{1}{3}\overline{\psi_L}\gamma_\mu\psi_L + \frac{4}{3}\overline{t_R}\gamma_\mu t_R - \frac{2}{3}\overline{b_R}\gamma_\mu b_R\right]^2 \tag{16}$$

where the λ^a are color matrices and $g_{3h} \gg g_{3\ell}, g_{1h} \gg g_{1\ell}$. The κ_3 terms are uniformly attractive; were they alone, they would generate large m_t **and** m_b. The κ_1 terms, in contrast, include a repulsive component for b. As a result, the combined effective interactions[14, 15, 40, 41, 42, 43, 44, 45, 46, 47, 48, 58, 59, 60, 61, 62]

$$\kappa^t = \kappa_3 + \frac{1}{3}\kappa_1 > \kappa_c > \kappa_3 - \frac{1}{6}\kappa_1 = \kappa^b \tag{17}$$

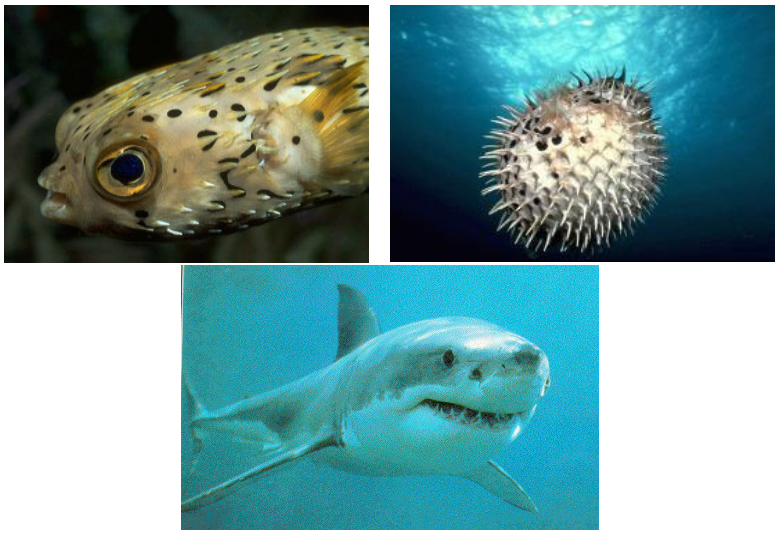

Figure 10: (a) Puffer fish. (b) Inflated spiny puffer fish. (c) Great white shark.

can be super-critical for top, causing $\langle \bar{t}t \rangle \neq 0$ and a large m_t, and sub-critical for bottom, leaving $\langle \bar{b}b \rangle = 0$.

The benefits of including new strong dynamics for the top quark are clear in TC2 models.[15, 58, 59, 60, 61, 62] Because technicolor is responsible for most of electroweak symmetry breaking, $\Delta\rho \approx 0$. Direct contributions to $\Delta\rho$ are avoided because the top condensate provides only $f \sim 60$ GeV; indirect contributions are not an issue if the technifermion hypercharges preserve weak isospin. The top condensate yields a large top mass. ETC dynamics at $M_{ETC} \gg 1$ TeV generate the light m_f without large FCNC and contribute only ~ 1 GeV to the heavy quark masses so there is no large shift in R_b.

Three classes of models of new strong top dynamics with distinctive spectra are known as topcolor [14, 40], flavor-universal extended color [15, 62, 63], and top seesaw [54]. Exotic particles in these models include colored gauge bosons (topgluons, colorons), color-singlet gauge bosons (Z'), composite scalar states (top-pions, q-pions), and heavy fermions (usually, but not always [19, 20], weak singlets). Because strong top dynamics is the subject of a talk by B. Dobrescu in these proceedings, it suffices here to note briefly that numerous searches for these new states have been attempted or proposed. For example, CDF has searched [64] for topgluons and Z' in heavy quark final states, and the potential for finding $Z' \to \tau\tau \to e\mu$ at Run II has been discussed in [21]. Limits on flavor-universal colorons are summarized in ref. [65]. The phenomenology of the new weak-singlet quarks present in top seesaw models has been discussed in refs. [66, 67]. Searches for the composite scalars of strong top dynamics models are analogous to the widely discussed methods for finding the extra scalars in multiple-Higgs models.

5. Summary

Dynamical symmetry breaking models, such as extended technicolor, use familiar gauge dynamics to give form to that most elusive quarry (fig. 10a), the origin of mass. Modern dynamical theories such as non-commuting ETC, low-scale technicolor, or topcolor-assisted technicolor have evolved in response to the pressures applied by increasingly precise measurements of observables such as R_b, flavor-changing neutral currents, and $\Delta\rho$. All of these theories offer intriguing and distinctive signatures (fig. 10b), many of which are discussed in these Proceedings by M. Narain. As the next round of collider experiments (fig. 10c) begins, I wish them "Good Hunting!"

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