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Mechanisms of Electroweak Symmetry Breaking The Role of a Heavy Top Quark

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Abstract

The dynamical basis of electroweak symmetry breaking remains an outstanding puzzle of elementary particle physics research. I review various mechanisms proposed to explain the origins of electroweak symmetry breaking with a particular focus on the possible role of the heavy top quark recently discovered at Fermilab.

Introduction.

Dynamical models based on local gauge symmetries provide the basis for our present understanding of the interactions of elementary particles. However, the masses of all known elementary particles appear to violate the local electroweak symmetries. Various mechanisms have been proposed to dynamically break these symmetries and generate the masses of all the quarks, leptons and gauge bosons. The top quark is the most massive of the known elementary particles and may play a special role.

Models of electroweak symmetry breaking can be separated into two classes. The first class contains models where the interactions associated with electroweak symmetry breaking are weakly interacting at the symmetry breaking scale of a few hundred GeV/c^2. These models include the Standard Model, the supersymmetric standard model and the top quark condensate models. At very high energy scales these models may become strongly interacting. For the other class of models, the dynamics becomes strongly interacting at scales close to the electroweak scale. These models include technicolor models, topcolor models, visible sector supersymmetry breaking and strong WW dynamics.

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Standard Higgs Model.

The Standard Model introduces an elementary Higgs scalar multiplet, an electroweak doublet field, in addition to the observed quarks, leptons and gauge vector bosons. The Higgs multiplet has Yukawa couplings to the fermions, gauge couplings to the vector bosons, and self-interactions. With a negative effective mass, the Higgs field develops a vacuum expectation value different from zero which breaks the electroweak symmetries. Particle masses are generated via the couplings to the Higgs field. The Standard Model has the following Higgs structure,

$$\begin{aligned} \mathrm{L} = &\overline{\mathrm{Q}}\mathrm{G}_u U \cdot H + \overline{Q} G_d D \cdot H^c + \overline{L} G_\ell E \cdot H^c \\ &+ \left(D^\mu H^+ D_\mu H\right) - \lambda\left(H^+ H\right)^2 + m^2\left(H^+ H\right) \end{aligned} \tag{1}$$

where G_u, G_d, G_ℓ are the coupling matrices for the up-quarks, down-quarks and leptons, respectively. The gauge boson dynamics is $SU(3)_{color} \otimes (SU(2) \otimes U(1))_{ew}$. The Standard Model has the remarkable feature of naturally protected masses and a natural Cabbibo-Kobayashi-Maskawa structure which suppresses flavor-changing neutral currents for the three generations of fermions.

The Standard Model could provide a precise description of the low energy dynamics. It may also be valid at scales far above the electroweak symmetry breaking scale. In this case the effective coupling constants are known to evolve, and we must have a consistent description of the physics at every scale that the model applies. The coupling constant evolution is described by the scale dependence of the coupling matrices, G_u, G_d, G_ℓ and may be computed by analyzing loop corrections. At lowest order, the top quark Yukawa coupling evolves according to the equation,

$$\mu\partial\mu g_t = g_t \cdot \left(\tfrac{9}{2} g_t^2 - 8 g_c^2 - \tfrac{9}{4} g^2 - \tfrac{17}{12} g'^2\right) / 16\pi^2 \tag{2}$$

where g_t is the diagonalized top quark Yukawa coupling and g_c, g, g' are the gauge couplings. This evolution is shown in Figure 1. It has the remarkable property that the physical mass of the top quark can not be too large if the model is to apply over a large range of scales, and the effective Yukawa coupling is not to "blow up". This behavior implies a pseudo-fixed-point structure for the evolution equations [1]. In Figure 2, the combined renormalization flow of the top quark and Higgs coupling constants are shown. In the Standard model, the fixed point value of the top quark mass is about 220 GeV/c^2 if the model is to remain valid to scales of order 10^{15} - 10^{19} GeV.

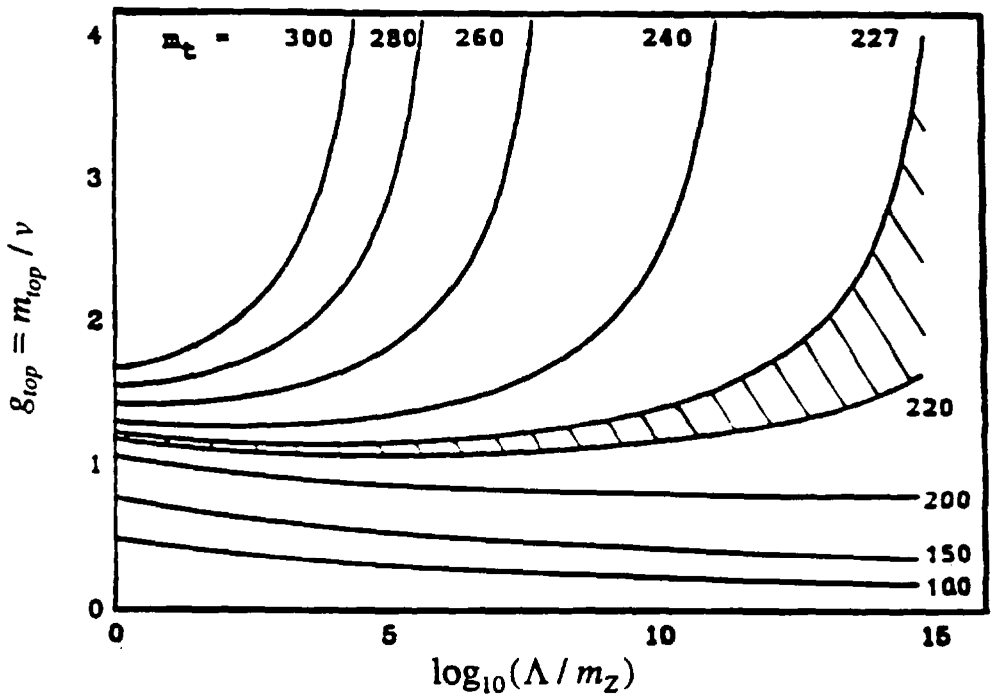

Figure 1. Top quark Yukawa coupling constant evolution in the Standard Model showing the pseudo-fixed point structure. From M. Lindner, in "Heavy Flavors", A. Buras and M. Lindner, eds., World Scientific, Singapore (1992), pg. 693.

However, the Standard Model is known to have a hierarchy problem for the Higgs mass which is not naturally protected as are all of the other particle masses. If we use a momentum cutoff to define the effective field theory at a given scale, then the Higgs mass will vary quadratically with the cutoff,

$$\Delta m^2 \to -4\sum_f m_f^2 \cdot \Lambda_f^2 + (2m_W^2 + m_Z^2 + m_h^2) \cdot \Lambda_b^2 \tag{3}$$

This behavior constitutes a fine tuning problem if the electroweak scale is generated within a large hierarchy [2]. When viewed from a high energy scale, the 'bare' Higgs mass would have to be carefully adjusted to yield the known scale of electroweak symmetry breaking. For many this represents a fatal flaw in the Standard model or implies that the Standard Model can only apply over a very limited range of energy scales near the electroweak scale.

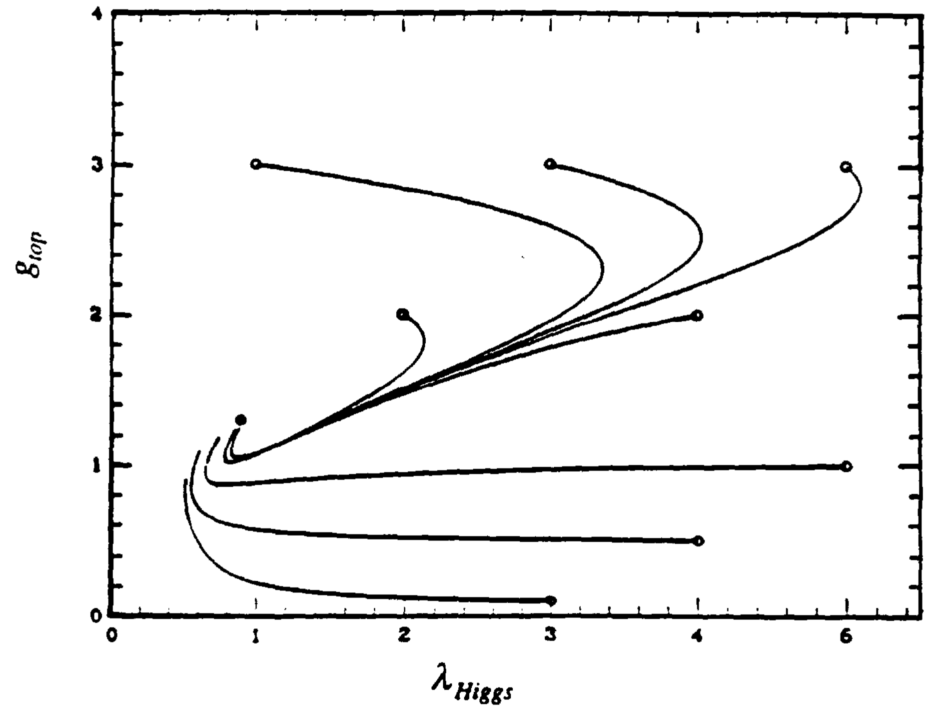

Figure 2. Renormalization flow of top quark and Higgs couplings in the Standard Model. From W. Bardeen and C. Hill, in "Heavy Flavors", A. Buras and M. Lindner, eds., World Scientific, Singapore (1992), pg. 649.

In perturbation theory, the quadratic divergence is an artifact of regularization. The standard model has a classical scale invariance which is broken by logarithms due to the running couplings. If the approximate scale invariance is maintained, then the quadratic divergence only reflects an explicit breaking of the scale invariance which should not be part of the perturbative expansion. Nonperturbative effects could still generate a hierarchy problem. Dynamical symmetries can relate the fermion cutoff, Λ_f, and the boson cutoff, Λ_b, in a natural way canceling the quadratic dependence and avoiding the hierarchy problem. Supersymmetry is an example of a dynamical symmetry of this type.

Supersymmetry.

In supersymmetric models, the fermi-bose symmetry protects the electroweak hierarchy. The minimal supersymmetric standard model, MSSM, requires the introduction of two Higgs fields in chiral supermultiplets. The Yukawa couplings are a direct generalization from the Standard Model,

$$L = QG_uU^c \cdot H_2 + QG_dD^c \cdot H_1 + LG_\ell E^c \cdot H_1 + \mu H_1 \cdot H_2 \tag{4}$$

$$+ \text{ soft supersymmetry breaking terms.}$$

Even in the presence of soft SUSY breaking terms, the Higgs mass terms are not renormalized by large perturbative corrections and the electroweak hierarchy can be preserved. The electroweak scale is then related to the scale of soft supersymmetry breaking. The explanation of the electroweak hierarchy becomes a question of a hierarchy that determines the scale of supersymmetry breaking.

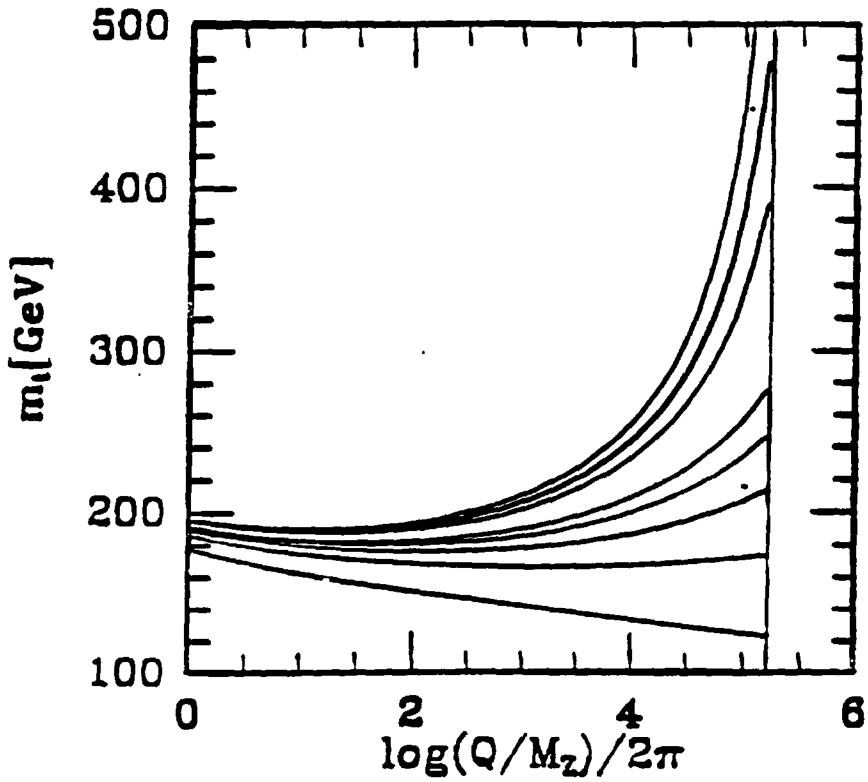

Figure 3. Top quark Yukawa coupling constant evolution in a minimal supersymmetric extension of the Standard Model showing the fixed point structure. From W. Bardeen, M. Carena, S. Pokorski and C. Wagner, Phys. Lett. B320, 110(1994).

As in the case of the Standard Model, the Yukawa couplings of the MSSM evolve when the theory is viewed at different reference scales. However, the supersymmetry modifies the evolution of the couplings. In the MSSM, the evolution equation for the top quark coupling becomes

$$\mu\partial\mu g_t = g_t \cdot \left(6g_t^2 - \tfrac{16}{3}g_c^2 - 3g^2 - \tfrac{13}{9}g'^2\right) / 16\pi^2 \qquad (5)$$

and is shown in Figure 3. As in the Standard Model, the MSSM predicts an infrared pseudo-fixed-point [3] for the evolution for the top coupling constant. The pseudo-fixed-point solution implies

$$m_{top} \to (190-200)\ GeV \cdot \sin\beta, \qquad \tan\beta \le 30 \qquad (6)$$

using the two-loop evolution equations. For appropriate values of $\tan\beta$, the predicted top quark mass falls in the range observed by CDF and DØ at Fermilab.

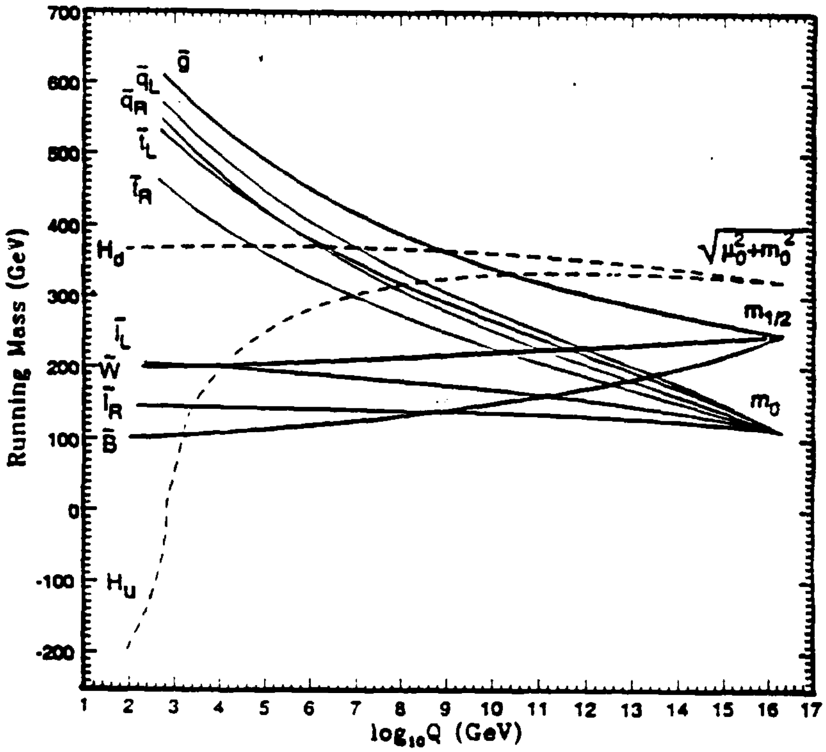

Figure 4. Scalar mass evolution in constrained MSSM showing radiative electroweak symmetry breaking due to the heavy top quark. From G. Kane, C. Kolda, L. Roszkowski and J. Wells, Phys. Rev. D49, 6173(1994).

A heavy top quark may also play a crucial role in generating electroweak symmetry breaking. It is usually assumed that the soft SUSY breaking terms generate a common scalar mass for the chiral matter fields at high energy. The matter fields include the supermultiplets for quarks, leptons and the two Higgs fields. The soft breaking terms are subject to radiative corrections and the scalar masses evolve when viewed at lower energy scales. The large top quark Yukawa coupling drives the up-Higgs field mass terms negative which triggers the electroweak symmetry breaking [4]. The running of scalar mass terms is shown in Figure 4.

The heavy top quark also plays an important role in supersymmetric GUT models. In many of these models, the bottom quark and the tau lepton have a common Yukawa coupling constant at the GUT scale. Below the GUT scale, these couplings evolve separately producing the considerable difference in the masses generated by the electroweak symmetry breaking as observed at low energy [5]. The evolution is governed by

$$\mu\partial\mu(m_b / m_\tau) = (m_b / m_\tau)\cdot\left(g_t^2 - \tfrac{16}{3}g_c^2 + \cdots\right) / 16\pi^2 \tag{7}$$

and is sensitive to the strong coupling constant, g_c^2 and the top quark Yukawa coupling constant, g_t^2. Hence, the bottom/tau mass ratio is uniquely sensitive to the top quark Yukawa coupling at high scales and, therefore, to its fixed point structure. For values of the bottom quark mass in the range, $M_b \sim 4.9\text{-}5.2$ GeV/c^2, and large values of the strong coupling constant, $\alpha_s \sim 0.125$, as required by the gauge coupling unification, the evolution causes the ratio to undershoot its unification value. Large values of the top quark Yukawa coupling would be required to compensate driving the top quark mass within 10% of its fixed point value. Large threshold effects may be required to achieve agreement with the full GUT picture and $b-\tau$ unification.

Supersymmetric models have many different aspects which fit together in intriguing ways. The heavy top quark is seen to play a crucial role in making different pieces fit into a unified picture relevant to all known energy scales.

Top Quark Condensates.

If there are no elementary Higgs particles, electroweak symmetry breaking could be generated by attractive short range interactions among the known particles. If the attraction is sufficiently strong, then a vacuum instability can occur, as in BCS theory, and the stable vacuum state is one with a dynamically broken symmetry. Because of its heavy mass, the top quark is a natural focus for this dynamics and forms the basis of top quark condensate models [6].

The short range dynamics of the top quark can be modeled by an attractive four fermion interaction, as in the Nambu-Jona-Lasinio model, NJL. This interaction must preserve the gauge symmetries of the Standard Model since the electroweak interactions are only broken dynamically. In the simplest version, only the top quark is important for the dynamics and the interaction has the form,

$$L = G \cdot (\overline{Q}_L t_R) \cdot (\bar{t}_R Q_L) \tag{8}$$

In the broken symmetry vacuum state, the top quark becomes massive and a chiral condensate of top quark pairs forms. Higg's particles are composite $t-\bar{t}$ states, and the effective theory at low energies is identical with the usual Standard Model except that the top quark mass is pushed close to its fixed point value.

The condensate breaks the electroweak symmetries, and the scale of electroweak symmetry breaking can be computed in terms of the top quark mass. In the NJL approximation, the electroweak scale is given by

$$v^2 \approx (3/16\pi^2) \cdot m_{top}^2 \cdot \log\left(\Lambda^2 / m_{top}^2\right) \qquad \text{(NJL)} \tag{9}$$

where Λ is composite scale and effective cutoff of the four-fermion dynamics in Eq.(8). Renormalization group methods can be used to improve the accuracy of the NJL predictions. To provide sufficient electroweak symmetry breaking, the top quark must be very heavy although limited by its fixed point value. If there is a large hierarchy between the electroweak scale and the composite scale, then the fixed point predictions are very stable,

$$m_{top} \approx 220GeV \text{ (full minimal model)}, \quad \approx 165\ GeV \text{ (NJL)} \tag{10}$$

assuming a composite scale, $\Lambda \approx 10^{15} GeV$. The top quark condensate model also makes predictions for the mass of the physical Higgs particle

$$m_{higgs} \approx 1.2 \cdot m_{top} \text{ (full minimal model)}, \quad \approx 2 \cdot m_{top} \text{ (NJL)} \tag{11}$$

The recent measurements of the top quark mass at Fermilab would appear to rule out the minimal version of the top quark condensate model.

The model also requires a fine tuning of the coupling constant in Eq.(8) to produce a large hierarchy. This fine tuning could have a dynamical explanation; the four-fermi interactions could be generated by a separate dynamically broken gauge dynamics at high energy, and the effective NJL coupling, corresponding to massive gauge boson exchange, would be determined dynamically by a gap equation that could require a large hierarchy for stability. Another possibility is a fourth generation with heavier quarks and leptons and a massive neutrino which would contribute to the electroweak symmetry breaking in addition to the top quark. In this case, the composite scale could be much lower, and the extensive fine tuning could be avoided. It is also possible to construct supersymmetric versions of the top quark condensate model [7] where the fixed point predictions of the top quark mass in Eq.(10) are lowered to the range observed by the Fermilab experiments. The supersymmetric models still have difficulties in understanding the composite scale and motivating composite picture for the Higgs multiplets.

Technicolor models.

Strong dynamics at a scale close to the electroweak symmetry breaking scale is an attractive alternative to models which remain perturbative up to very high energy. As in quantum chromodynamics, the scale where the dynamics becomes strongly interacting could then be determined from the logarithmic running of a coupling constant rather than an ad hoc scale of electroweak or supersymmetry breaking.

Technicolor models [8] are based on a strong, asymptotically free gauge field theory. The full gauge group is normally thought to be a direct product of the technicolor group and the gauge group of the Standard Model,

$$SU(N_{TC}) \otimes (SU(3) \otimes SU(2) \otimes U(1))_{SM} \qquad (12)$$

Technifermions are introduced which carry both technicolor and electroweak dynamics. Technifermions condense when the technicolor interactions become strong and dynamically break the electroweak symmetries and generate masses for the electroweak gauge bosons, m_W and m_Z,

$$\langle \overline{T}^a T_b \rangle \neq 0, \qquad \text{Technicolor scale} \sim 1\ \text{TeV} \qquad (13)$$

Dynamical electroweak symmetry breaking at low energy provides a solution to the hierarchy problem. A new strong spectrum (technirhos, etc) with strong dynamics is predicted. Strong production of new states (techniquarks) and resonance production of normal particles (top quarks) is possible. "Realistic" family structures for the technifermions generally imply a large set of global chiral symmetries which are dynamically broken by the strong technicolor dynamics and produce a spectrum of pseudo-Goldstone bosons which are light technifermion bound states.

Fermion masses for quarks and leptons must be generated through mixing with technifermions, the extended technicolor (ETC) interactions [9]. ETC dynamics is usually associated with the dynamical breaking of a larger unified gauge dynamics whose representations involve both technifermions, T, and normal fermions, F. At low energies, the ETC dynamics has the form,

$$L_{ETC} = G_{FF}(\bar{F}F)^2 + G_{TF}(\bar{T}T)\cdot(\bar{F}F) + G_{TT}(\bar{T}T)^2 \tag{14}$$

$$G \approx 1/\Lambda_{ETC}^2 \approx g_{ETC}^2/m_{ETC}^2$$

G_{FF} is constrained by bounds on flavor-changing neutral currents (FCNC) which imply $\Lambda_{ETC} > 200-1000\ TeV$. G_{TF} generates the quark and lepton masses,

$$m_F \approx G_{TF}\langle\bar{T}T\rangle \tag{15}$$

G_{TT} generates contributions to the pseudo-Goldstone boson masses,

$$f_P^2 \cdot m_P^2 \approx G_{TT}\left\langle [Q^P,[Q^P,(\bar{T}T)^2]]\right\rangle + \text{gauge contributions} \tag{16}$$

The technicolor models have many problems. The FCNC constraints limit the strength of the ETC interactions and make it difficult to generate sufficiently large masses for the quarks, leptons and pseudo-Goldstone bosons. Here, the heavy top quark is a particular embarrassment. The lack of a natural GIM mechanism in technicolor models produces an unnatural CKM structure. Precision electroweak measurements place strong constraints on the technicolor contributions to the S, T, U parameters which are sensitive to physics beyond the Standard Model. These constraints place severe limits on technicolor model building, and no "minimal standard model" exists for technicolor dynamics.

Since technicolor models are based on theories with strong dynamics, it is difficult to make firm predictions. Originally it was assumed that the strong technicolor dynamics could be scaled from the low energy dynamics of quantum chromodynamics. In QCD the dynamics is localized at the QCD scale which would imply technicondensates behave as

$$\langle\bar{T}T\rangle \approx \Lambda_{TC}^3, \qquad \Lambda_{TC} \text{ - technicolor scale} \tag{17}$$

The spectrum of states could also be scaled from QCD. With this scaling property, the quark, lepton and pseudo-Goldstone masses are much too small. These predictions killed the original formulations of technicolor dynamics.

Large numbers of technifermions can slow the running of the technicolor gauge coupling constant and produce "walking" technicolor models [10]. The slow running of the coupling constant over a large energy range can generate large effective anomalous dimensions for the technicolor operators and enhance the size of the technicondensates,

$$\langle \bar{T}T \rangle \approx \Lambda_{TC}^{3} \cdot \left(\Lambda_{ETC} / \Lambda_{TC}\right)^{\gamma}, \quad \gamma \text{ - anomalous dimension} \tag{18}$$

with $\Lambda_{ETC} >> \Lambda_{TC}$, $\gamma > 0$. This effect enhances the generation of masses for both fermions and pseudo-Goldstone bosons. This enhancement may not be sufficient for the third generation: top, bottom and tau.

Additional enhancements can be achieved if some of the four-fermi ETC interactions have near critical NJL dynamics [11],

$$G_{TT} \to G_{critical} \tag{19}$$

This mechanism is similar to the top quark condensate dynamics without the associated fine-tuning problems. The technifermion condensates are enhanced even if the NJL dynamics is subcritical due to the near formation of technifermion bound-states.

The top quark mass may require special treatment and models with several scales of technicolor dynamics, perhaps associated with the generation structure [12]. Also top-bottom and top-tau mass splittings may be difficult to achieve from ETC dynamics alone. Near critical instabilities of the NJL dynamics may be required to enhance the intra-family hierarchies.

Despite the severe dynamical constraints, semi-realistic technicolor model-building is possible [13]. There exist examples of extended technicolor models which address the problems of FCNC, CP-violation, quark and lepton mass hierarchies and precision electroweak measurements. There have also been many papers which focus on critical tests of technicolor dynamics and technicolor model building including the constraints of precision electroweak measurements (S,T,U) [14], B-meson decays $(B \to s\gamma,\ B \to \mu^+\mu^- X)$ [15], anomalous gauge boson couplings $(Zb\bar{b},\ WWg,\ WWZ\ \ etc)$ [16] and pseudo-Goldstone boson phenomenology [17].

Topcolor dynamics.

As the most massive elementary particle yet known, the top quark may play as special role in the dynamics of electroweak symmetry breaking. Topcolor models formulate a new strong dynamics involving specifically the third generation of quarks and leptons [18]. This dynamics enhances the scale of top quark condensates and provides a natural explanation of the large top quark mass. If these models prove consistent, there will be a rich and testable phenomenology near or below the scale of electroweak symmetry breaking.

Topcolor dynamics is based on a new strong dynamics involving the third generation at high energies. A separate, weak dynamics involves the first two generations. Near the 1 TeV scale, these two interactions combine to produce the usual gauge dynamics of the Standard Model. In this scheme, the gauge dynamics has the structure,

$$G_3 \otimes G_{12} \to G_{SM}, \qquad \text{broken near the 1 TeV scale} \qquad (20)$$

where G_3 couples to the top quark or third generation with a strong coupling,

$$\frac{1}{g_{SM}^2} = \frac{1}{h_3^2} + \frac{1}{h_{12}^2}, \qquad h_3 / h_2 \equiv \cot\theta >> 1 \qquad (21)$$

and G_{12} couples to the first two generations. Topcolor gauge symmetry breaking produces massive gauge bosons, topgluons, whose exchange produces massive gauge bosons whose exchange generates a strong attractive dynamics involving the top quark (third generation). This dynamics triggers condensates which break the electroweak symmetries and generate the top quark (third generation) mass, etc.

A particular example of topcolor dynamics has been studied by C. Hill [19] where the gauge dynamics is specified by

$$G_3 = (SU(3) \otimes U(1))_{\text{third generation, color and hypercharge}}$$

$$G_{12} = (SU(3) \otimes U(1))_{\text{first and second generations}} \qquad (22)$$

$$G_{SM} = SU(3)_c \otimes U(1)_Y$$

After the symmetry breaking of Eq.(20), the effective action for the new dynamics is

$$L' = -\frac{4\pi\kappa}{M_B^2}\left[\bar{t}\gamma_\mu \frac{\lambda^A}{2} t + \bar{b}\gamma_\mu \frac{\lambda^A}{2} b\right]^2$$

$$- \frac{4\pi\kappa_Y}{M_{Z'}^2}\left[\frac{1}{3}\bar{\Psi}_L\gamma_\mu\Psi_L + \frac{4}{3}\bar{t}_R\gamma_\mu \frac{\lambda^A}{2} t_R - \frac{2}{3}\bar{b}_R\gamma_\mu \frac{\lambda^A}{2} b_R\right]^2 \tag{23}$$

where $\kappa = g_C^2 \cot^2\theta / 4\pi$, $\kappa_Y = g'^2 \cot^2\theta'$. The induced strong topcolor interaction is an attractive, isospin symmetric gauge interaction for top and bottom quarks. In addition, the induced U(1) interactions are attractive in the top quark channel and repulsive in the bottom quark channel. The model presumes that the top quark interactions produce a strong top quark condensate, $\langle\bar{t}t\rangle$, but that the interactions are subcritical in the bottom quark channel, and no bottom quark condensate is formed. The large top quark mass results from the top quark condensate. The bottom quark does not get its mass directly from condensate but from instanton effects of the strong gauge dynamics. The model predicts the existence of pseudo-Goldstone bosons, top-pions, associated with the symmetry breaking.

Various models based on topcolor dynamics have been constructed. Electroweak symmetry breaking may be dominated by the topcolor dynamics [18] with all masses being related to a version of top quark condensation. In alternative schemes, topcolor is used in combination with additional technicolor dynamics to produce the full electroweak symmetry breaking [20].

Topcolor models are expected to have some common phenomenological implications. The pseudo-Goldstone bosons, the top-pions, have a low dynamical scale and a strong coupling,

$$f_{TP} \approx 50\ \ GeV,\ \ g_{TP} \approx m_{top} / \sqrt{2} f_{TP} \tag{24}$$

The third generation has new strong dynamics. The new gauge dynamics could be observed directly as resonant production in the $t\bar{t}$ invariant mass distributions as shown in Figure 5 for topgluon production and in Figure 6 for the Z' of the U(1) dynamics. Anomalous couplings for the top and bottom quark and the tau lepton should also be generated. The enhancement of R_b observed at LEP may be a first indication of this new dynamics. A number of processes including $b \to s\gamma$, $\Delta S = 2$, $\Delta C = 2$, are sensitive to topcolor dynamics. The ρ-parameter is also sensitive to particular aspects of topcolor models. These processes already place strong constraints on topcolor model-building [21].

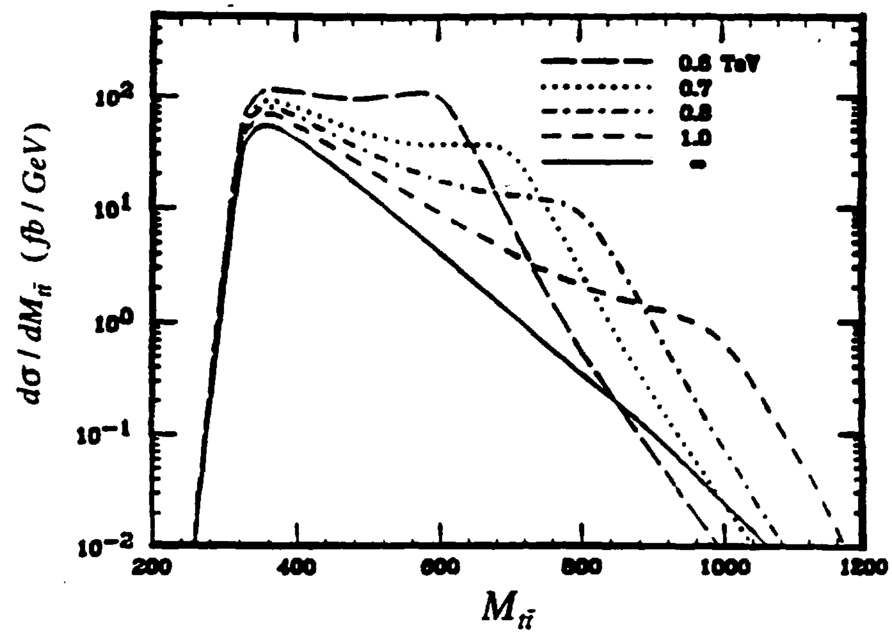

Figure 5. Structure in the $t\bar{t}$ mass distribution at the Tevatron for various topgluon masses, $\Gamma_{tg} \approx 0.2 \cdot M_{tg}$. From C. Hill and S. Parke, Phys. Rev. D49, 4454(1994).

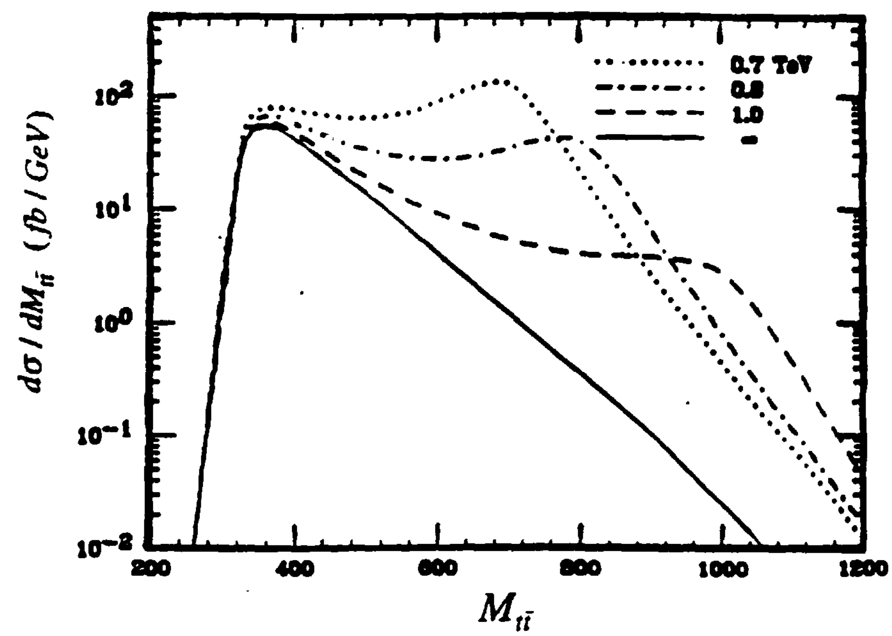

Figure 6. Structure in the $t\bar{t}$ mass distribution at the Tevatron for various Z' masses, $\Gamma_{Z'} \approx 0.2 \cdot M_{Z'}$. From C. Hill and S. Parke, Phys. Rev. D49, 4454(1994).

Visible Sector SUSY Breaking.

In the usual MSSM, supersymmetry breaking is thought to arise through the interactions of supergravity with a "hidden" sector where SUSY is broken dynamically. Electroweak symmetry breaking occurs through radiative corrections driven by the large Yukawa couplings of the heavy top quark. Supersymmetry breaking could also be generated by visible-sector gauge boson dynamics. Here the symmetry breaking can occur at lower energy scales and perhaps avoid some of the naturalness and cosmological problems of the hidden-sector supergravity models.

Visible-sector models involve a supersymmetric gauge sector with dynamical breaking of supersymmetry via instantons. The symmetry must be transmitted to the conventional SUSY sector via "messenger" fields. This can be achieved by gauging an additional U(1) charge, "messenger hypercharge". Radiative electroweak symmetry breaking can be achieved, depending on the SUSY breaking scales, via loop effects involving the large top quark Yukawa coupling. The gravitino in the LSP with a mass determined by the SUSY breaking scale, M_s, and the Planck scale, M_P, with

$$m_{3/2} = M_s^2 / M_p \tag{25}$$

$m_{3/2}$ is constrained by cosmology to be less than 10 KeV and implies that M_s must be less than about 10^7 GeV [22]. These models dynamically break a U(1) R-symmetry generating a pseudo-Goldstone boson, the R-axion, with a mass generated by Planck scale physics. The lower limit on this mass implies that the SUSY breaking scale must be larger that about 10^5 GeV [23].

Strong WW dynamics.

The absence of a low energy manifestation of electroweak symmetry breaking implies a strong dynamics for the vector bosons at a scale of about 1 TeV. This scales is determined by the strong interaction of the longitudinal components of the vector bosons. A compact description of these interactions can be achieved through a gauge invariant, chiral Lagrangian whose Goldstone bosons become the longitudinal components of the W and Z gauge bosons [24]. The effective action could also describe the effective dynamics of technicolor and topcolor models.

A model independent analysis can be made through a systematic momentum expansion of the chiral Lagrangian [25]. There are two dimension-two operators which generate corrections to the W and Z boson masses. There are eleven CP-conserving, dimension-four operators and three CP-violating, dimension-four operators. Of these sixteen operators, only four operators contribute directly to the two-point functions for the gauge bosons. These operators generate the usual electroweak scale, the rho parameter and the S and U parameters, the oblique corrections [26]. Other operators contribute to vertex and higher point functions. Many papers estimate the magnitude of the various terms in particular models.

Phenomenological implications of strong WW scattering can be studied by using the effective chiral Lagrangian expansion. Detailed studies of signals and backgrounds for probing the strong WW sector have been made for a variety of accelerators and detectors, particularly the SSC and the LHC [27].

Conclusions.

The discovery of the top quark and the determination of its large mass, 180±12 GeV, has important implications for understanding the mechanisms of electroweak symmetry breaking. If electroweak symmetry breaking is due to an elementary Higgs multiplet with a large hierarchy to new physics beyond the electroweak scale, then the pseudo-fixed-point structure of both the Standard Model and its supersymmetric extension may play an essential role in determining the low energy dynamics. In the supersymmetric model, the fixed-point structure has a direct impact on the top quark (and possibly the bottom quark) Yukawa coupling as well as the mechanism of radiative electroweak symmetry breaking and the predictions for $b-\tau$ or $t-b-\tau$ unification in SUSY GUT models. The predictions of the top quark condensate models also focus on the fixed-point structure of the effective standard model dynamics observed at low energies.

The heavy top quark also has a direct impact on models where the Higgs mechanism is due to composite structure at the electroweak scale. Technicolor and ETC scenarios require an understanding of nonperturbative strong dynamics. Models have evolved from the QCD based analysis which could not support a large top quark mass to Walking and NJL enhanced models which may be able to generate the large masses observed for the third generation fermions. Topcolor models focus on special aspects of the strong dynamics associated directly with the third generation. The heavy top quark would signal a new strong dynamics for the third generation which must be visible at the electroweak scale.

The present phenomenological success of the minimal Standard Model does not directly point to the physical mechanism of electroweak symmetry breaking. This success already places strong constraints on the directions of new physics associated with these mechanisms. As we begin to probe the physics associated with the scale of top quark, we may begin to unravel the puzzle of electroweak symmetry breaking. Novel mechanisms may emerge as we confront the new physics to be observed by present and future experiments.

In the near future, an upgraded Tevatron will provide new opportunities to directly probe the physics of the top quark and this new scale of dynamics. LEP II will provide precise information on the W-boson. Preliminary evidence for new physics associated with supersymmetry or composite structure may begin to emerge from these facilities.

In the future the LHC will provide a unique probe of the broad spectrum of physics possibilities which must become visible at the electroweak scale. New linear colliders may allow us to probe these new scales in e^+e^- and $\gamma\gamma$ collisions. Further upgrades to the Tevatron may allow for a high intensity quark-antiquark collider. Muon colliders may also provide a new avenue to higher energy probes of the electroweak physics and beyond.

WITH THE DISCOVERY OF THE TOP QUARK, WE BEGIN A NEW ERA!

References.

[1] B. Pendleton and G. Ross, Phys. Lett. 98B, 291(1981); C. Hill, Phys. Rev., D24, 691(1981); E. Paschos, Z. Phys. C26, 235(1984); C. Hill, C. Leung and S. Rao, Nucl. Phys. B262, 517(1985).

[2] S. Weinberg, Phys. Lett. 82B, 387(1979); G. 't Hooft, Lectures at the 1979 Cargese Summer Institute (1979), pg. 135.

[3] J. Bagger, S. Dimopoulos and E. Massó, Phys. Rev. Lett. 55, 920(1985).

[4] L. Ibanéz and G. Ross, Phys. Lett. 110B, 215(1982); K. Inoue, A. Kakuto, H. Komatsu, and S. Takeshita, Prog. Theor. Phys. 68, 927(1982); erratum, ibid, 70, 330(1983); L. Alvarez-Gaumé, J. Polchinsky and M. Wise, Nucl. Phys. B221, 495(1983); J. Ellis, J. Hagelin, D. Nanopolous, and K. Tamvakas, Phys. Lett. 125B, 275(1983); M. Carena, M. Olechowski, S. Pokorski and C. Wagner, Nucl. Phys. B419, 213(1994).

[5] H. Arason, D. Castaño, B. Keszthelyi, S. Mikaelian, E. Piard, P. Ramond and B. Wright, Phys. Rev. D47, 232(1993); V. Barger, M. Berger, P. Ohmann and R. Phillips, Phys. Lett. B314, 351(1993); P. Langacker and N. Polonsky, Phys. Rev. D49, 1454(1994); W. Bardeen, M. Carena, S. Pokorski and C. Wagner, Phys. Lett. B320, 110(1994).

[6] Y. Nambu, NagoyaWorkshop on Strong Dynamics (1988); EFI Preprint 89/08 (1989); V. Miransky, M. Tanabashi and K. Yamawaki, Phys. Lett. B221, 177(1989); Mod. Phys. Lett. A4, 1043(1989); W. Bardeen, C. Hill and M. Lindner, Phys. Rev. D41, 1647(1990); A. Blumhofer, R. Dawid and M. Lindner, Phys. Lett. B360, 123(1995)

[7] M. Carena, T. Clark, C. Wagner, W. Bardeen and K. Sasaki, Nucl. Phys. B369, 33(1992).

[8] S. Weinberg, Phys. Rev. D13, 974(1976); D19, 1277(1979); L. Susskind, Phys. Rev. D20, 2619(1979); S. Dimopoulos and L. Susskind, Nucl. Phys. B155, 237(1979).

[9] E. Eichten and K. Lane, Phys. Lett. 90B, 125(1980).

[10] B. Holdom, Phys. Lett. 150B, 301(1985); T. Appelquist, D. Karabali and L. Wijewardhana, Phys. Rev. Lett. 57, 957(1986); K. Yamawaki, M. Bando and K. Matumoto, Phys. Rev. Lett. 56, 1335(1986); T. Appelquist and L. Wijewardhana, Phys. Rev. D35, 774(1987).

[11] T. Appelquist, M. Einhorn, T. Takeuchi and L. Wijewardhana, Phys. Lett. B220, 223(1989); V. Miransky and K. Yamawaki, Mod. Phys. Lett. A4, 129(1989); K. Matumoto, Prog. Theor. Phys. Lett. 81, 277(1898); V. Miransky, M. Tanabashi and K. Yamawaki, Phys. Lett. B221, 177(1989).

[12] B. Holdom, Phys. Rev. Lett. 60, 1233(1988); T. Appelquist and O. Shapira, Phys. Lett. B249, 327(1990).

[13] T. Appelquist and J. Terning, Phys. Rev. D50, 2116(1994); N. Evans, Preprints SWAT-33 (1994); YCTP-P12-94 (1994); S. Chivukula, M. Dugan, M. Golden and E. Simmons, Preprint BUHEP-95-09 (1995).

[14] S. Chivukula, M. Dugan and M. Golden, Phys. Lett. B292, 435(1992); Phys. Rev. D47, 2930(1993); R. Sundrum and S. Hsu, Nucl. Phys. B391, 127(1993), T. Appelquist and J. Terning, Phys. Rev. D47, 3075(1993); C. Hill, D. Kennedy, T. Onogi and H-L. Yu, Phys. Rev. D47, 2940(1993); S. Chivukula, B. Dobrescu and J. Terning, Preprint BUHEP-95-22 (1995).

[15] L. Randall and R. Sundrum, Phys. Lett. B312, 148(1993); B. Grinstein, Y. Nir and J. Soares, Phys. Rev. D48, 3960(1993); C. Carone, E. Simmons and Y. Su, Phys. Lett. B344, 287(1995).

[16] U. Mahanta, Phys. Rev. D46, 3123(1992); R. Peccei, S. Peris and X. Zhang, Nucl. Phys. B349, 305(1991); T. Appelquist and G.-H. Wu, Phys. Rev. D48, 3235(1993); S. Chivukula, E. Gates, E. Simmons and J. Terning, Phys. Lett. B311, 157(1993); S. Chivukula, E. Simmons and J. Terning, Phys. Lett. B346, 284(1995); C. Carone, E. Simmons and Y. Su, Phys. Lett. B344, 287(1995); C. Hill and X.-M. Zhang, Phys. Rev. D51, 3563(1995); S. Chivukula, Preprint BUHEP-95-17 (1995); C.-X. Yue, Y.-P. Kuang, G.-R. Lu and L.-D. Wan, Preprint TUIMP-TH-95/62 (1995).

[17] T. Appelquist and G. Triantaphyllou, Phys. Rev. Lett. 69, 2750(1992); E. Eichten and K. Lane, Phys. Lett. B327, 129(1994); S. Chivukula, R. Rosenfeld, E. Simmons and J. Terning, Preprint BUHEP-95-07 (1995); K. Lane, Phys. Lett. B357, 624(1995).

[18] C. Hill, Phys. Lett. B266, 419(1991); R. Bönish, Phys. Lett. B268, 394(1991); S. Martin, Phys. Rev. D46, 2197(1992); Phys. Rev. D45, 4283(1992); M. Lindner and D. Ross, Nucl. Phys. B370, 30(1992); C. Hill, D. Kennedy, T. Onogi and H.-L. Yu, Phys Rev. D47, 2940(1993); B. Holdom, Phys. Lett. B336, 85(1994)

[19] C. Hill, Phys. Lett. B345, 483(1995).

[20] C. Hill, Phys. Lett. B345, 483(1995); K. Lane and E. Eichten, Phys. Lett. B352, 382(1995); N. Evans, Preprint SWAT-33 (1994); YCTP-P12-94 (1994); K. Lane, Phys. Lett. B357, 624(1995).

[21] S. Chivukula, B. Dobrescu and J. Terning, Phys. Lett. B353, 289(1995).

[22] M. Dine and A. Nelson, Phys. Rev. D48, 1277(1993); M. Dine, A. Nelson and Y. Shirman, Phys. Rev. D51, 1362(1995).

[23] J. Bagger, E. Poppitz and L. Randall, Nucl. Phys. B426, 3(1994).

[24] J. Cornwall, D. Levin and G. Tiktopoulos, Phys. Rev. D10, 1145 (1974); M. Chanowitz and M.-K. Gaillard, Nucl. Phys. B261, 379(1985).

[25] T. Appelquist and C. Bernard, Phys. Rev. D22, 200(1980); A. Longhitano, Phys. Rev. D22, 1166(1980); Nucl. Phys. B188, 118(1980); T. Appelquist and G.-H. Wu, Phys. Rev. D48, 3235(1993); Phys. Rev. D51, 240(1995); H.-J. He, Y.-P. Kuang and C.-P. Yuan, Phys. Rev. D51, 6463(1995).

[26] M. Peskin and T. Takeuchi, Phys. Rev. Lett. 65, 964(1990); Phys. Rev. D46, 381(1992).

[27] M. Chanowitz, in "Perspectives on Higgs Physics", G. Kane, ed. (1992); J. Bagger, V. Barger, K. Cheung, J. Gunion, T. Han, G. Ladinsky, R. Rosenfeld and C.-P. Yuan, Fermilab-PUB-93/040-T (1993); M. Chanowitz and W. Kilgore, Phys. Lett. B322, 147(1994); Phys. Lett. B347, 387(1995).