

Electroweak Symmetry Breaking via a Top Condensation Seesaw Mechanism

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We propose a new mechanism in which all of the electroweak symmetry breaking is driven by a natural top quark condensate, produced by “top-color” interactions at a multi-TeV scale. The scheme allows the observed top quark mass, and acceptable T and S parameters, by invoking a seesaw mechanism involving mixing of the top with an additional isosinglet quark. [S0031-9007(98)07205-6]

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The proximity of the measured top quark mass, m_t , to the electroweak scale hints that electroweak symmetry breaking (EWSB) has its origin in dynamics associated with the top quark. An explicit realization of this idea is the top condensation mechanism [1], in which the top-antitop quark pair acquires a vacuum expectation value, much like the chiral condensate of QCD or the electron pairing condensate of BCS superconductivity. The $\langle \bar{t}t \rangle$ condensate has the correct quantum numbers for EWSB, and it connects the electroweak scale directly to m_t . However, the original schemes of this type [1,2] are problematic: (i) they require the scale of new physics to be very large, of the order of the grand unified theory (GUT) scale, $\sim 10^{15}$ GeV; (ii) they require an unnatural cancellation of large quadratic mass scales; (iii) they lead in the standard model to the renormalization group fixed point prediction $m_t \sim 230$ GeV which is too high [3] (note though that, in the minimal supersymmetric standard model, the fixed point prediction is reasonable [4]).

Subsequently, a specific gauge dynamics, “top-color,” was proposed to generate a top condensate at the TeV scale [5]. This model views the heaviness of the top quark as a dynamical phenomenon which is essentially independent of EWSB, provides a specific model of the new dynamics, and attempts to place the scale of the dynamics at ~ 1 TeV to avoid fine-tuning issues. It produces associated pseudo-Nambu-Goldstone bosons with decay constant $f_{t\pi}$. Using the large- N , fermion loop approximation, one estimates

$$f_{t\pi}^2 \approx \frac{3}{4\pi^2} m_t^2 \ln\left(\frac{M}{m_t}\right). \quad (1)$$

With $M \sim 1$ TeV, we obtain $f_{t\pi} \approx 64$ GeV (EWSB, would require $f_{t\pi} = 246$ GeV). Hence, top-color has always been combined with additional dynamics, such as technicolor [6]. Indeed, this provides a solution to the serious problem of the large top quark mass within the context of technicolor models, and interesting top-color-assisted technicolor models (or TC²) have been constructed [7]. These models, though potentially viable, are somewhat cumbersome, and implications of limits on custodial symmetry violation [8] and other phenomenological

constraints [9,10] require dynamical assumptions which are difficult to analyze.

In the present Letter we propose a new mechanism, in which the EWSB occurs via the condensation of the top quark in the presence of an extra vectorlike, weak-isosinglet quark. The mass scale of the condensate is large, of the order of 0.6 TeV corresponding to the electroweak scale $v \approx 246$ GeV. The vectorlike isosinglet then naturally admits a seesaw mechanism, yielding the physical top quark mass, which is then adjusted to the experimental value. The choice of a natural multi-TeV scale for the top-color dynamics then determines the mass of the vectorlike quark.

There are several attractive features of this mechanism: (i) The model is relatively specific, the left-handed top quark being unambiguously identified as the electroweak-gauged condensate fermion, and the scheme uniquely specifies top-color, together with some new U(1) interactions, as the primary new strong interaction; (ii) the scheme is economical, requiring no additional weak isodoublets, and therefore easily satisfies the constraints upon the S parameter, using estimates made in the large- N approximation; (iii) the constraint on custodial symmetry violation, i.e., the value of the $\Delta\rho_* \equiv \alpha T$ parameter, is easily satisfied, being principally the usual m_t contribution suppressed by the squared ratio of the mixings between the top and the vectorlike quark. We mention that we were led to these schemes by considering fermionic bound states in the context of strong dynamical models. We will not elaborate this aspect presently, but note that the vectorlike quark in our model may potentially arise from within the dynamics without being introduced *ad hoc* (an example of this type is given in Ref. [11]).

Consider the embedding of the standard model into a top-color scheme: $SU(3)_1 \times SU(3)_2 \times SU(2)_W \times U(1)_{1_1} \times U(1)_{2_1} \times U(1)_{B-L}$ gauge group where the $SU(3)_1 \times U(1)_{1_1}$ [$SU(3)_2 \times U(1)_{2_1}$] acts only on the third (first and second) generation quarks and leptons. The $U(1)_{B-L}$ charges are $x < 0$ for t_R , $1/3$ for the other quarks, and -1 for leptons (including right-handed neutrinos). In addition to the observed quarks and leptons we

include a new fermion, χ , whose left- and right-handed components transform identically to t_R , except that the $U(1)_{B-L}$ charge of χ_L is x , such that the anomalies cancel. At a multi-TeV scale the gauge symmetries are broken as follows: (i) $SU(3)_1 \times SU(3)_2 \rightarrow SU(3)_{\text{QCD}}$, leaving a degenerate octet of massive ‘‘colorons’’; (ii) $U(1)_1 \times U(1)_2 \rightarrow U(1)_Y$, and (iii) $U(1)_{B-L}$ is broken, leaving two heavy gauge bosons, Z_1 and Z_{B-L} ; (iv) $SU(2)_W$ is unbroken.

We emphasize that this specific choice of the gauge groups and particle content is considered presently for the purpose of illustration. To achieve the desired physics in this scheme will require some fine-tuning of the gauge coupling constants to a few percent of critical values. We believe that this fine-tuning can be alleviated in more general schemes, and in fact, we expect to generalize the mechanism when it is extended to give the light quark and lepton masses. Moreover, we will work in the fermion-bubble approximation to leading order in N_c , the number of colors. This is a crude approximation to the dynamics, so the specific estimates made here are expected to be reliable to roughly only a factor of 2 or so.

Integrating out the colorons yields an effective Lagrangian for the third generation fields:

$$\mathcal{L}_{\text{eff}}^c = -\frac{4\pi\kappa}{M^2} \left(\bar{q}\gamma_\mu \frac{\lambda^A}{2} q + \bar{\chi}\gamma^\mu \frac{\lambda^A}{2} \chi \right)^2, \quad (2)$$

where $q = (t, b)$, the coloron mass is M , $\kappa = (g_s^2 \cot^2 \theta)/8\pi$, where g_s is the QCD coupling, and θ is the usual $SU(3)_1 \times SU(3)_2$ mixing angle [5,6] (the first and second generations feel effects that are proportional to $\tan^2 \theta$ and mixing with the third generation that are of order 1); we assume $\cot \theta \gg 1$. For critical κ the effects of $\mathcal{L}_{\text{eff}}^c$ alone would produce an $SU(3)$ symmetric condensate with $\langle \bar{t}t \rangle = \langle \bar{b}b \rangle = \langle \bar{\chi}\chi \rangle$, and an $SU(3)$ octet of Nambu-Goldstone bosons (NGB's).

Integrating out the two massive $U(1)$ bosons yields the effective Lagrangians:

$$\mathcal{L}_{\text{eff}}^1 = -\frac{4\pi\kappa_1}{M_1^2} \left(\frac{1}{3} \bar{q}_L \gamma_\mu q_L + \frac{4}{3} \bar{t}_R \gamma_\mu t_R - \frac{2}{3} \bar{b}_R \gamma_\mu b_R + \frac{4}{3} \bar{\chi} \gamma^\mu \chi \right)^2, \quad (3)$$

and

$$\mathcal{L}_{\text{eff}}^{B-L} = -\frac{4\pi\kappa_{B-L}}{M_{B-L}^2} \left(\frac{1}{3} \bar{q}_L \gamma_\mu q_L + x \bar{t}_R \gamma_\mu t_R + \frac{1}{3} \bar{b}_R \gamma_\mu b_R + x \bar{\chi}_L \gamma^\mu \chi_L + \frac{1}{3} \bar{\chi}_R \gamma^\mu \chi_R \right)^2, \quad (4)$$

as well as four-fermion operators involving the third generation leptons. Here M_1 and M_{B-L} are the Z_1 and

Z_{B-L} masses, $\kappa_1 = g'^2 \cot^2 \theta'/(8\pi)$, g' is the hypercharge gauge coupling, θ' is the mixing angle between the $U(1)_1 \times U(1)_2$ gauge bosons, $\kappa_{B-L} = g_{B-L}^2/(8\pi)$, g_{B-L} is the $U(1)_{B-L}$ gauge coupling. We shall assume, for simplicity, that all of the massive gauge bosons have a common mass M . These interactions are typically attractive and nonconfining. In addition, electroweak preserving mass terms are allowed in the low energy theory:

$$\mathcal{L}_{\text{mass}} = -\mu_{\chi\chi} \bar{\chi}_L \chi_R - \mu_{\chi t} \bar{\chi}_L t_R. \quad (5)$$

Note that the $\bar{\chi}_L \chi_R$ mass term is not $U(1)_{B-L}$ invariant, but it is induced below the scale M if the χ couples to the $U(1)_{B-L}$ breaking vacuum expectation value (VEV).

In the presence of the full effects of $\mathcal{L}_{\text{eff}}^1$ and $\mathcal{L}_{\text{eff}}^{B-L}$, the theory can produce a different pattern of condensates, i.e., a different pattern of chiral symmetry breaking, than in the case of the pure $\mathcal{L}_{\text{eff}}^c$. To see this we note that the mass-gap equations for the model take the generic form

$$m_{AB} = z_{AB} m_{AB} F_{AB}(m_{\chi\chi}, m_{\chi t}, m_{t\chi}, m_{tt}, \dots), \quad (6)$$

where the indices A and B stand for χ, t , and b , m_{AB} is the dynamical mass generated by the condensate $\langle \bar{A}_L B_R \rangle$, and the functions F_{AB} depend on all nine dynamical masses. The coefficients z are combinations of the $SU(3)_1 \times U(1)_1 \times U(1)_{B-L}$ gauge couplings and charges:

$$z_{AB} = \frac{2}{\pi} \left(\frac{4}{3} \kappa + Y_A Y_B \kappa_1 + X_A X_B \kappa_{B-L} \right), \quad (7)$$

where Y and X are the $U(1)_1$ and $U(1)_{B-L}$ charges, respectively. Our charge assignment gives the following inequalities:

$$\begin{aligned} z_{t\chi} &> z_{tt}, z_{tb} > z_{\chi b}, \\ z_{\chi t} &> z_{\chi\chi} > z_{tt}. \end{aligned} \quad (8)$$

A necessary condition for having a nonzero dynamical mass m_{AB} is that at least one of the three $z_{A/B}$ coefficients is above a critical value, $z_{\text{crit}} = 1$, or that the corresponding current mass μ_{AB} is nonzero. In the low energy effective theory this condition corresponds to the requirement of having a negative squared mass or a tadpole term for the composite scalar formed in the $\bar{A}_L B_R$ channel [12]. We need the formation of the $\langle \bar{\chi}_L t_R \rangle$ and $\langle \bar{t}_L \chi_R \rangle$ condensates, and therefore we require $z_{\chi t}, z_{t\chi} > 1$. Choosing the three gauge couplings such that $z_{tb} < 1$ ensures that the $\bar{q}_L b_R$ and $\bar{\chi}_L b_R$ channels are subcritical, so that b_R does not participate in condensates. Furthermore, if z_{tt} is subcritical, the $\bar{q}_L t_R$ condensate will be aligned with $\bar{q}_L \chi_R$. In addition, we assume for convenience that $z_{\chi\chi} < 1$. All of the above conditions on the z coefficients are satisfied provided

$$\max\left(\frac{4\kappa_1}{3} + \frac{x\kappa_{B-L}}{4}, -\frac{\kappa_1}{6} + \frac{\kappa_{B-L}}{12}\right) < \frac{3\pi}{8} - \kappa < \min\left(\frac{\kappa_1}{3} + \frac{\kappa_{B-L}}{12}, \frac{4\kappa_1}{3} + \frac{3x^2}{4}\kappa_{B-L}\right). \quad (9)$$

In this case, the most general dynamical mass matrix [up to an $SU(2)_W$ transformation] is given in the weak eigenstate basis by

$$\mathcal{L} = -(\bar{t}_L, \bar{\chi}_L) \begin{pmatrix} m_{tt} & m_{t\chi} \\ m_{\chi t} & m_{\chi\chi} \end{pmatrix} \begin{pmatrix} t_R \\ \chi_R \end{pmatrix} + \text{H.c.} \quad (10)$$

Therefore, we need to solve the subset of four coupled gap equations included in Eq. (6), where the functions F can be computed by keeping the weak eigenstates in the external lines, and the χ and t mass eigenstates running in the loop. Keeping only the quadratic and logarithmic divergences, with a physical cutoff given by the mass of the heavy gauge bosons, we obtain

$$\begin{aligned} F_{\chi\chi} &= 1 - \left(m_{t\chi}^2 + m_{\chi t}^2 + m_{\chi\chi}^2 + \frac{m_{tt}m_{t\chi}m_{\chi t}}{m_{\chi\chi}} \right) \\ &\quad \times \frac{1}{M^2} \ln \left(\frac{M^2}{m_{\chi\chi}^2} \right) + \frac{\mu_{\chi\chi}}{z_{\chi\chi}m_{\chi\chi}}, \\ F_{\chi t} &= 1 - \left(m_{tt}^2 + m_{\chi t}^2 + m_{\chi\chi}^2 + \frac{m_{tt}m_{t\chi}m_{\chi\chi}}{m_{\chi t}} \right) \\ &\quad \times \frac{1}{M^2} \ln \left(\frac{M^2}{m_{\chi\chi}^2} \right) + \frac{\mu_{\chi t}}{z_{\chi t}m_{\chi t}}, \\ F_{t\chi} &= 1 - \left(m_{tt}^2 + m_{t\chi}^2 + m_{\chi\chi}^2 + \frac{m_{tt}m_{\chi t}m_{\chi\chi}}{m_{t\chi}} \right) \\ &\quad \times \frac{1}{M^2} \ln \left(\frac{M^2}{m_{\chi\chi}^2} \right), \\ F_{tt} &= 1 - \left(m_{tt}^2 + m_{t\chi}^2 + m_{\chi t}^2 + \frac{m_{t\chi}m_{\chi t}m_{\chi\chi}}{m_{tt}} \right) \\ &\quad \times \frac{1}{M^2} \ln \left(\frac{M^2}{m_{\chi\chi}^2} \right). \end{aligned} \quad (11)$$

In what follows we will be interested for simplicity in the case where the $\langle \bar{\chi}_L \chi_R \rangle$ condensate is significantly larger than the others, such that the seesaw condition,

$$m_{\chi\chi}^2 \gg m_{\chi t}^2, m_{t\chi}^2 \gg \frac{m_{\chi\chi}^2}{m_{\chi t}^2} m_{tt}^2, \quad (12)$$

is satisfied. A nontrivial solution to the set of gap equations can be easily found in this case (we assume $\mu_{\chi t} \ll \mu_{\chi\chi}$):

$$\begin{aligned} m_{\chi\chi} &\approx \frac{\mu_{\chi\chi}}{z_{t\chi} - z_{\chi\chi}} > 0, \\ m_{\chi t, t\chi}^2 &\approx \frac{M^2}{\ln(M^2/m_{\chi\chi}^2)} \left(1 - \frac{1}{z_{\chi t, t\chi}} \right) - m_{\chi\chi}^2 > 0, \quad (13) \\ m_{tt} &\approx -\frac{z_{tt}m_{\chi\chi}m_{\chi t}m_{t\chi}}{(1 - z_{tt})M^2} \ln \left(\frac{M^2}{m_{\chi\chi}^2} \right) < 0. \end{aligned}$$

The seesaw condition (12) requires a partial cancellation of the terms in the above expression for $m_{\chi t, t\chi}^2$, and a constraint on $\mu_{\chi\chi}$,

$$\frac{\mu_{\chi\chi}}{M^2} \ll (1 - z_{tt})(z_{t\chi} - z_{\chi\chi})^2. \quad (14)$$

It should be emphasized that it is possible to satisfy conditions (9) and (12) without excessive fine-tuning. To see this, note that the $SU(3)_1$ interactions are likely to be stronger than the $U(1)$'s, so that Eq. (9) gives the criticality condition

$$\kappa = \frac{3\pi}{8} \left[1 + O \left(\frac{\kappa_1}{\kappa}, \frac{\kappa_{B-L}}{12\kappa} \right) \right]. \quad (15)$$

According to the discussion in Ref. [8], the ratios κ_1/κ and $\kappa_{B-L}/(12\kappa)$ can be interpreted as the amount of fine-tuning in κ required by Eq. (9). From this point of view, we consider values of order 0.1–0.01 for these ratios to be acceptable. We return to an example of the numerical values of coupling constants after obtaining some phenomenological constraints on the masses.

As a consequence of condition (12), the physical top mass is suppressed by a seesaw mechanism:

$$m_t \approx \frac{m_{\chi t}m_{t\chi}}{m_{\chi\chi}} \left[1 + O(m_{\chi t, t\chi}^2/m_{\chi\chi}^2) \right]. \quad (16)$$

The electroweak symmetry is broken by the $m_{t\chi}$ dynamical mass. Therefore, the electroweak scale is estimated to be given by

$$v^2 \approx \frac{3}{4\pi^2} m_{t\chi}^2 \ln \left(\frac{M}{m_{\chi\chi}} \right). \quad (17)$$

It is easy to verify that this estimate is correct to first order in $(m_{t\chi}/m_{\chi\chi})^2$, by diagonalizing the mass matrix in Eq. (10) and computing the one-loop leading contribution to the W mass. This can also be seen in an effective Lagrangian analysis; this term comes from a dynamically generated composite Higgs boson with a very strong coupling $g \sim 4\pi/\sqrt{3} \ln(M/m)$ to $\bar{t}\chi$, with the usual VEV, where Eq. (17) is just the usual $m = gv/2$ “Goldberger-Treiman” relation. Thus, $v \approx 246$ GeV requires a dynamical mass

$$m_{t\chi} \approx 0.6 \text{ TeV} \quad (18)$$

for $M/m_{\chi\chi} \sim O(10)$ [values of $M/m_{\chi\chi}$ much larger than $O(10)$ would need fine-tuning in the gap equation for $m_{\chi\chi}$, while values below ~ 5 would need the inclusion of the next-to-leading order terms in the gap equations]. From Eq. (16) it follows that a top mass of 175 GeV requires

$$\frac{m_{\chi t}}{m_{\chi\chi}} \approx 0.3. \quad (19)$$

One might worry at this stage that the four-fermion interactions (3) and (4), which are custodial- $SU(2)$ violating, inducing the top-bottom mass splitting, may also lead to a large contribution to the electroweak parameter T (equivalently, $\Delta\rho_*$). T can be estimated in the fermion-bubble large- N approximation as

$$T \approx \frac{3m_t^2}{16\pi^2\alpha(M_Z^2)v^2} \frac{m_{t\chi}^2}{m_{\chi t}^2} \left[1 + O(m_{\chi t, t\chi}^2/m_{\chi\chi}^2) \right], \quad (20)$$

where α is the fine structure constant. The two-loop corrections to T involving the operators (2), which are significant for $M \sim 1$ TeV as discussed in Ref. [8], are small in our case due to the large M [compared to the result in Ref. [8], there is a suppression of $(1 \text{ TeV}/M)^2 \lesssim 10^{-3}$, and an enhancement of $(m_{t\chi}/m_t)^4 \sim 100$ from χ loops]. Moreover, the S parameter is small, the leading contributions coming from a χ - t loop diagram, of order $(1/2\pi)(m_{\chi t}/m_{\chi\chi})$. The fit to the electroweak data of the S and T parameters in the standard model with a Higgs mass of 300 GeV gives [13] a 1σ ellipse in the S - T plane whose projection on the T axis is $T = 0.03 \pm 0.34$. Requiring that our model does not exceed the 1σ upper bound on T , we obtain the following constraint:

$$\frac{m_{t\chi}}{m_{\chi t}} \leq 0.55. \quad (21)$$

The mass ratios (19) and (21) can arise naturally. For example, the set of values $m_{t\chi} \approx 0.5$ TeV, $m_{\chi t} \approx 0.9$ TeV, $m_{\chi\chi} \approx 2$ TeV, $M \sim 50$ – 70 TeV would require $\kappa_1 \approx 0.02$, $\kappa_{B-L} \approx 0.3$, and $1 - 8\kappa/(3\pi) \approx 0.02$, for $x = -1/6$ and $\mu_{\chi\chi} \approx 30$ GeV. Note that the vectorlike quark has a mass approximately equal to $m_{\chi\chi}$, and therefore the phenomenological constraints on its mass [14] are satisfied. A consequence of the mixing between the t and χ quark eigenstates is that the V_{tb} element of the Cabibbo-Kobayashi-Maskawa matrix is given by the cosine of the mixing angle. Therefore, the decrease in $|V_{tb}|$ compared to the standard model value is given by $m_{t\chi}^2/2m_{\chi\chi}^2 \sim 1\%$, which is potentially relevant for the single top production at the run 3 of the Tevatron.

This model also implies the existence of pseudo-NGB's. The chiral symmetry of the four-fermion interactions (2)–(4) is $SU(2)_W \times U(1)_{\chi_L} \times U(1)_{\chi_R} \times U(1)_{t_R} \times U(1)_{b_R}$, where the four $U(1)$'s are the chiral symmetries of χ_L , χ_R , t_R , and b_R . One of the linear combinations of the $U(1)$ generators corresponds to the hypercharge gauge group $U(1)_Y$, while the remaining combinations are generators of global $U(1)$'s. The condensates $\langle \bar{\chi}_L \chi_R \rangle$ and $\langle \bar{\chi}_L t_R \rangle$ break $U(1)_{\chi_L} \times U(1)_{\chi_R} \times U(1)_{t_R} \times U(1)_{b_R}$ down to $U(1)_Y \times U(1)_{b_R}$ giving rise to two neutral NGB's which are linear combinations of $\bar{\chi} \gamma_5 \chi$, $\bar{\chi} \gamma_5 t$, and $\bar{t} \gamma_5 t$. These NGB's receive masses of order

$$M_{\text{NGB}} \approx \sqrt{\mu_{\chi A} m_{\chi A}}, \quad (22)$$

where A stands for χ or t , due to the explicit breaking of the global $U(1)$'s by the Dirac mass terms in Eq. (5). Note that the lower bound on the mass of a neutral pseudo-NGB that couples only to the top quark is much below the bound on the Higgs boson mass [15], so that even a current mass $\mu_{\chi t}$ of order 1 GeV is sufficient.

In summary, we have presented a simple mechanism of electroweak symmetry breaking based on the dynamics of the top quark with a seesaw mechanism involving a

vectorlike quark. The model produces an acceptable top-quark mass, is consistent with other electroweak data, and does not require excessive fine-tuning. It remains for one to understand the necessary extension of the scheme to give masses and mixing angles to all quarks and leptons, and to construct attractive schemes for top-color breaking.

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