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Electroweak Symmetry Breaking as a Consequence of Compact Dimensions

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Abstract

It has been shown recently that the Higgs doublet may be composite, with the left-handed top-bottom doublet and a new right-handed anti-quark as constituents bound by some four-quark operators with nonperturbative coefficients. I show that these operators are naturally induced if there are extra space dimensions with a compactification scale in the multi-TeV range. The Higgs compositness is due mainly to the Kaluza-Klein modes of the gluons, while flavor symmetry breaking may be provided by various fields propagating in the compact dimensions. I comment briefly on the embedding of this scenario in string theory.

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Recently it has been shown that a bound state formed of the left-handed top-bottom quark doublet and a new right-handed anti-quark can play viably the role of the Standard Model Higgs doublet [1, 2]. The only ingredients necessary in this framework are some four-quark interactions, suppressed by a multi-TeV scale M_1. One is then led to ask what is the origin of these non-renormalizable interactions. The traditional answer is that they are produced by the exchange of some heavy gauge bosons associated with the breaking of a larger gauge group down to the QCD group. The simplest choice of this sort is the topcolor [3] scheme: $SU(N_c)_1 \times SU(N_c)_2 \to SU(N_c)_{\rm C}$, where $N_c = 3$ is the number of colors, the third generation quarks are charged under $SU(N_c)_1$, and the first two generations are charged under $SU(N_c)_2$. This choice has several nice features. For example, the asymptotic freedom of the topcolor gauge group allows the heavy gauge bosons to be rather strongly coupled at the scale M_1 without problems with a Landau pole. Also, it is technically convenient because it allows the use of the large N_c expansion.

The drawback is that one has to include additional structures that would induce the spontaneously breaking of topcolor. Moreover, this embedding of the $SU(N_c)_{\rm C}$ color group may be seen as somewhat artificial. The embedding can be improved by allowing all quarks to transform non-trivially only under $SU(N_c)_1$ [4] and then to embed this group in a gauged flavor or family symmetry [5, 6]. However, such an extension appears to require a more complicated topcolor breaking sector.

It is therefore legitimate to ask the question: is it possible to avoid the extension of the gauge group while inducing the desired four-quark operators and retaining the nice features of topcolor ? In this letter I point out that the required four-quark operators are naturally induced if the Standard Model gauge fields propagate not only in the usual 4-dimensional Minkowski space-time but also in a compact manifold.

The first step in showing this is to recall that the existence of the extra dimensions is manifested within the 4-dimensional space-time through a tower of Kaluza-Klein (KK) modes associated with each of the fields propagating in the compact manifold (see *e.g.* [7, 8]). The KK modes of a massless gauge boson have masses between $M_1 = 1/R_{\rm max}$, where $R_{\rm max}$ is the largest compactification radius, and the fundamental scale associated with strongly coupled quantum gravity, M_*. The fermions couple only to the vector excitations of the gauge bosons, and the couplings are identical with those of the gauge bosons. Therefore, the KK excitations of the gluons give rise in the low energy theory to flavor universal four-quark operators:

$$\mathcal{L}_{\rm eff}^c = -\frac{c g_s^2}{2M_1^2}\left(\sum_q \overline{q}\gamma_\mu T^a q\right)^2 , \tag{1}$$

where q are all the colored fields, T^a are the $SU(N_c)_{\rm C}$ generators, and g_s is the QCD gauge coupling at the scale M_1. The dimensionless coefficient $c > 0$ sums the contributions of all gluonic KK modes, so that it depends on the number of extra dimensions and on the compactification radii.

Higher dimensional operators are also induced in the low energy theory, but the contribution from a heavy gluonic KK mode, of mass M_n, is suppressed compared to the contribution to c by powers of M_n/M_1. Therefore, the effects of the higher dimensional operators may be neglected in the effective theory. By contrast, in topcolor models the higher dimensional quark operators induced by the heavy gauge bosons are usually ignored for convenience; this procedure may be physically reasonable but so far has not been mathematically motivated.

The contact interaction $\mathcal{L}_{\rm eff}^c$ is attractive in the scalar channel, so that spinless $\overline{q}_L q_R$ bound states form. Their properties can be studied using an effective potential formalism. In the large N_c limit only the left-right current-current part of $\mathcal{L}_{\rm eff}^c$ contributes to the effective potential for the composite scalars. Note that a Fierz transformation of this current-current interaction gives the well known Nambu–Jona-Lasinio interaction. Hence, the large N_c limit and the Nambu–Jona-Lasinio model are equivalent approximations of the gluonic KK dynamics. For c larger than a critical value, the composite scalars acquire vevs and break the electroweak symmetry [9]. Ignoring the renormalization group evolution of c above the scale M_1, one can find the critical value to leading order in $1/N_c$:

$$c_{\rm crit} = \frac{2\pi}{N_c \alpha_s} , \tag{2}$$

where $\alpha_s = g_s^2/4\pi$ is the strong coupling constant at the scale M_1. The important feature here is that the chiral phase transition is second order as c is varied. This property is expected to remain true beyond the large N_c approximation [10]. In the absence of excessive fine-tuning, *i.e.* if c is not very close to the critical value, the electroweak scale of 246 GeV indicates that M_1 is in the multi-TeV range, or smaller.

It would be useful to find the constraint on the dimensionality and topology of the compact manifold necessary for electroweak symmetry breaking to occur. For simplicity, consider a δ-dimensional torus with constant radii R_l, $l = 1, ..., \delta$. The spectrum of the

KK excitations of a massless field is given by

$$M^2_{n_1,\ldots,n_\delta} = \sum_{l=1}^{\delta} \frac{n_l^2}{R_l^2} \ , \tag{3}$$

where n_l are integers (KK excitation numbers). In what follows, the KK mass levels will be denoted by M_n with n integer ($1 \le n \le n_{\rm max}$ where $M_{n_{\rm max}} \approx M_*$), and their degeneracy by D_n. To compute the coefficient of the four-quark operator, one would need to integrate out the $D_{n_{\rm max}}$ modes of mass $M_{n_{\rm max}}$, then to use the renormalization group evolution for c from $M_{n_{\rm max}}$ down to $M_{n_{\rm max}-1}$, and to repeat these steps for each KK mass threshold until the lightest states are integrated out. This would give the coefficient c at the scale M_1 as a function of M_*/R_l and $g_s(M_1)$.

Fortunately it is not necessary to perform this computation in order to show that there are compact manifolds which induce electroweak symmetry breaking. Furthermore, a continuity argument shows that c may be tuned very close to the critical value. To see this, note that the contribution from any gluonic KK state to c is positive, so that truncating the tower of states at some $n_{\rm tr} < n_{\rm max}$ gives $c(n_{\rm tr}) < c(n_{\rm max}) = c$. Furthermore, the running of c may be ignored if $M_{n_{\rm tr}}$ is sufficiently close to M_1. For example, if the δ extra dimensions have the same compactification radius, R, then the truncation of the KK tower at $M_{n_{\rm tr}} = 2/R$ yields

$$c(n_{\rm tr} = 4) = \sum_{n=1}^{4} \frac{D_n}{n} \ , \tag{4}$$

where the degeneracy is given by

$$D_n = 2^n \, \delta! \left[\frac{\theta(\delta - n)}{n! \, (\delta - n)!} + \frac{\theta(n-4)\theta(\delta - n + 3)}{8 \, (n-4)! \, (\delta - n + 3)!} \right] \ , \tag{5}$$

with the step function $\theta(x \ge 0) = 1$. Note that this expression for the degeneracy is valid only for $n \le 7$.

For $\delta = 4$, and M_1 in the multi-TeV range [where $\alpha_s(M_1) \approx 0.08 - 0.06$, corresponding to $c_{\rm crit} \approx 26 - 35$],

$$c > c(4) \approx 35.7 > c_{\rm crit} \ , \tag{6}$$

which shows that four extra dimensions are sufficient for the composite Higgs doublets to acquire vevs. As a corollary, the Standard Model is not viable if there are 4 or more extra dimensions with a compactification scale below the fundamental Planck scale M_*, because the quarks would acquire dynamical masses of the order of the compactification scale.

Consider now the case $\delta = 1$. Ignoring the running of c, the sum over the contributions from all KK states yields

$$c = \sum_{n_1=1}^{n_{\max}} \frac{2}{n_1^2} < \frac{\pi^2}{3} < c_{\rm crit} \ . \tag{7}$$

Even if the contributions from the running of c happen to be positive, it seems highly unlikely that they amount to the factor of order 10 required to drive c over the critical value. Thus, if there is only one compact dimension, it is fair to conclude that the criticality condition is not satisfied, and the electroweak symmetry remains unbroken.

Coming back to the super-critical $\delta = 4$ case, one can decrease three of the four radii, recovering continuously the case of only one extra dimension. This is a second order phase transition from the phase in which electroweak symmetry is broken to the unbroken phase. The boundary of the phase transition corresponds to a set of tori with radii $R_1, ..., R_\delta$, where

$$2 \le \delta \le 4 \ . \tag{8}$$

A more precise estimate of c might raise the lower bound on δ. On the other hand, the above arguments show that there are 4-dimensional manifolds, with a hierarchy of compactification radii [at least one radius has to be shorter than $1/(2M_*)$], which yield c very close to the critical value. This result opens up the possibility of constructing realistic models of Higgs compositness based on compact dimensions.

Although the KK excitations of the gluon may induce electroweak symmetry breaking, they do not provide flavor symmetry breaking. For super-critical c, the effects of $\mathcal{L}^c_{\rm eff}$ alone would produce an $SU(N_f)$ symmetric condensate and an $SU(N_f)$ adjoint of Nambu–Goldstone bosons (where N_f is the number of quark flavors). All the quarks would acquire the same dynamical mass, related to the electroweak scale. It is therefore necessary to identify a source of flavor symmetry breaking. Also, as explained in ref. [1, 2], at least one new quark, χ, should be introduced such that a $\overline{t}_L \chi_R$ dynamical mass of order 0.5 TeV is induced, leading to the observed W and Z masses.

The KK excitations of the hypercharge gauge boson give rise to four-fermion operators which are attractive for the up-type quarks and repulsive for the down-type quarks:

$$\mathcal{L}^Y_{\rm eff} = -\frac{c g'^2}{M_1^2} \left(\frac{1}{3} \overline{\psi}^i_L \gamma_\mu \psi^i_L + \frac{4}{3} \overline{u}^i_R \gamma_\mu u^i_R - \frac{2}{3} \overline{d}^i_R \gamma_\mu d^i_R + \frac{4}{3} \overline{\chi} \gamma^\mu \chi \right)^2 \ , \tag{9}$$

where $i = 1, 2, 3$ is a generational index, $\psi^i_L = (u^i, d^i)_L$, and the lepton currents are not shown for simplicity. The vectorlike quark χ transforms under the Standard Model gauge group in the same representation as u^i_R. If the KK modes of the gluons yield c within about 10% of its critical value, then the combination $\mathcal{L}^c_{\rm eff} + \mathcal{L}^Y_{\rm eff}$ induces vevs only for the Higgs fields made up of the u, c, t and χ quarks.

At this stage it is necessary to introduce inter-generational flavor symmetry breaking. It is convenient to parametrize it using four-fermion operators of the following type:

$$\frac{\eta_{AB}}{M_1^2}\left(\overline{A}_L B_R\right)\left(\overline{B}_R A_L\right) \ , \tag{10}$$

with the notation $A_L = \psi^i_L, \chi_L$ and $B_R = u^k_R, d^i_R$, where $k = 1, 2, 3, 4$, and $u^4_R \equiv \chi_R$. Unlike the four-quark operators induced by the KK modes of the gluons, which are strongly coupled and give rise to deeply bound states, the four-fermion operators (10) can be treated perturbatively because the coefficients η_{AB} do not need to be larger than order one.

The origin of these operators can be found in different scenarios for the physics at scales above M_1. For example they can be produced by quantum gravitational effects provided the fundamental scale is of order M_1. This may occur in the truly strong coupling regime of string theory [11], as well as when there are large dimensions accessible only to gravitons [12]. A similar scenario is considered in ref. [13] for the purpose of providing flavor symmetry breaking in technicolor models.

An alternative would be the existence of scalars with positive squared masses of order M_1. If the fundamental scale M_* is much larger than M_1, their masses may be protected by supersymmetry. The exchange of such scalars would lead to the four-fermion operators (10) with coefficients η_{AB} dependent on the Yukawa couplings.

If one of the $\eta_{\psi^i u^k}$ coefficients, chosen by convention to be $\eta_{\psi^3\chi}$, is larger than the other eleven, then the condition

$$\eta_{\psi^3\chi} > \frac{2\pi^2}{N_c} - c\left(g_s^2 + \frac{8}{9N_c}g'^2\right) > \eta_{\psi^i u^k} \tag{11}$$

implies that only the $\overline{\chi}_R\psi^3_L$ Higgs doublet has a negative squared mass. As a result, the hierarchy between the top quark and the others will be naturally predicted.

In order to accommodate the observed masses of the W, Z and t, a few other conditions must be satisfied [1, 2]. First, the gauge invariant mass term $\mu_{\chi\chi}\overline{\chi}_L\chi_R$ has to be included, so that tadpole terms for the $\overline{\chi}_L\chi_R$ scalar are induced in the effective potential. For this purpose, the $\mu_{\chi\chi}$ mass parameter may be significantly smaller than the compactification scale M_1. Such a small mass may arise naturally, for example from the vev of a gauge singlet scalar which propagates in some compact dimensions which are inaccessible to χ

(a similar mechanism is used in ref. [15] to produce neutrino masses). Generically, the $\mu_{\chi u^i}\overline{\chi}_L u^i_R$ mass terms are also present.

A second condition for realizing the top condensation seesaw mechanism [1, 2] is to have the squared masses of the $\overline{\chi}_R\chi_L$ and $\overline{t}_R\chi_L$ scalars larger than the squared mass of the $\overline{\chi}_R\psi^3_L$ Higgs doublet, or equivalently:

$$\eta_{\psi^3\chi} > \eta_{\chi t}, \, \eta_{\chi\chi} \; . \tag{12}$$

These conditions ensure that the minimum of the effective potential for composite scalars corresponds to dynamical fermion masses only for the t and χ. The $\overline{t}_L\chi_R$ mass mixing, responsible for electroweak symmetry breaking, has to be of order 0.5 TeV, while the top mass measurements and the constraint on custodial symmetry violation require the fermion mass eigenvalues to be $m_t \approx 175$ GeV and $m_\chi \gtrsim 3$ TeV. In the absence of excessive fine-tuning, $m_\chi \sim \mathcal{O}(5)$ TeV corresponds to a compactification scale M_1 of order $10 - 100$ TeV.

The gluonic KK excitations are responsible for the existence of 49 composite complex scalars, each of the seven left-handed quark flavors binding to each of the seven right-handed ones (more composites may form if there are more χ quarks). Their mass degeneracy is lifted by the flavor non-universal four-fermion operators (10), but generically most of the physical states have masses of order m_χ [2]. However, the composite scalar spectrum includes three Nambu-Goldstone bosons which become the longitudinal W and Z, and a neutral Higgs boson which is always lighter than 1 TeV, and may be as light as $\mathcal{O}(100)$ GeV if the vacuum is close to the boundary of a second order phase transition [2]. In the decoupling limit, where $m_\chi \to \infty$, the low energy theory is precisely the Standard Model, with the possible addition of other composite states which may be light due to the vicinity of a phase transition.

So far, the electroweak symmetry is broken correctly, the top quark mass is naturally accommodated, and the χ quark has a mass of at least a few TeV. It remains to produce the masses and mixings of the other quarks and leptons. For this reason, consider the following four-fermion operators:

$$\frac{1}{M_1^2}\left(\overline{\chi}_R\psi^3_L\right)\left[\xi_{\psi^j u^k}\left(\overline{\psi}^j_L u^k_R\right) + \xi_{\psi^j d^i}\left(\overline{\psi}^j_L i\sigma_2 d^i_R\right) + \xi_{l^j\nu^i}\left(\overline{l}^j_L \nu^i_R\right) + \xi_{l^j e^i}\left(\overline{l}^j_L i\sigma_2 e^i_R\right)\right] \; , \tag{13}$$

where l^j_L, ν^i_R and e^i_R are the Standard Model lepton fields, and the dimensionless coefficients ξ are given by physics above M_1. These operators lead through the renormalization group evolution to Standard Model Yukawa interactions between the $\overline{\chi}_R\psi^3_L$ composite

Higgs doublet and the fermions. Therefore, in this picture the quark and lepton masses are produced by physics above the compactification scale. Another effect of the operators (13) is to mix the $\overline{\chi}_R \psi_L^3$ Higgs doublet with the other $\overline{q}_R^i \psi_L^j$ composite scalars.

The flavor breaking operators (10) and (13) do not produce large flavor-changing neutral current (FCNC) effects beyond those in the Standard Model. It remains to be shown however that these operators can be induced by some high energy physics without being accompanied by flavor-changing operators which produce large neutral-current effects. Another contribution to the FCNC's comes from the composite scalars, which have flavor breaking couplings to the quark mass-eigenstates. These are suppressed by Kobayashi-Maskawa elements, and they are not worrisome if the scalars are sufficiently heavy. It is beyond the scope of this letter to derive the constraints from FCNC's, but it is worth reiterating that the extra dimensions open up new ways of dealing with flavor. In addition, the low energy theory is also flexible: for example the four-quark operators may be kept flavor universal by allowing flavor violation to arise from the mass terms of an extended vectorlike quark sector [14, 5].

Note that the four-fermion operators induced in the four-dimensional flat space-time are independent of whether the fermions propagate or not in the extra δ dimensions. It is however striking that the new vectorlike quark, χ, required in the top condensation seesaw mechanism to account for the bulk of electroweak symmetry breaking, has the quantum numbers of the KK modes of the right-handed up-type quarks. The KK modes of the other standard fermions may be also present without affecting electroweak symmetry breaking. On the other hand, it appears more convenient to choose a tree level mass for χ below the scale of Higgs compositness [2], which may be hard to accommodate if χ is a KK excitation.

It is instructive to see how the scenario presented here may arise from string theory or M theory. The most convenient setting for studying large extra dimensions is within Type I string theory [8, 11, 16]. The closed string sector gives rise to the graviton and other neutral modes which propagate in the bulk of the 9+1 dimensional space-time. The open string sector gives rise to the gauge fields and the charged matter, which are restricted to propagate on a D9-brane or a D5-brane. Using T-duality transformations, one may obtain Type I$'$ string theories containing Dp-branes with $p \leq 9$. A $(\delta + 3)$-brane, with $2 \leq \delta \leq 4$, is necessary for containing the 3+1 dimensional flat space-time plus the δ compact dimensions that lead to a composite Higgs sector. Some of the δ compactification scales are expected to be of order $M_1 \sim 10 - 100$ TeV, while other must be slightly higher, in order to allow the coefficients of the four-quark operators to be close to criticality.

Due to the presence of the KK modes of the Standard Model gauge bosons, the gauge couplings tend to unify at a scale higher by about one order of magnitude than the compactification scale [8, 17]. The running of the gauge couplings in the model discussed here is different than in the Minimal Supersymmetric Standard Model because below the compactification scale there are no superpartners, there is a potentially complicated composite Higgs sector, and the heavy quark χ is also present. Nevertheless, the gauge coupling unification may be realized in various ways, due to the possible existence of additional states below or above the compactification scale. Alternately, the $SU(3)_C \times SU(2)_W \times U(1)_Y$ gauge groups may come from different types of branes, so that the gauge couplings need not unify precisely (the "unification" scale would be given in this case by the scale where the three gauge couplings come close to each other).

It is natural to identify the unification scale with the string scale: $M_s \sim \mathcal{O}(10^6)$ GeV. This prediction is a direct consequence of the fit to the electroweak scale and top-quark mass within the top seesaw theory of Higgs compositness. Once the string scale is specified, it is straightforward to determine the compactification radius, r', of the $6 - \delta$ dimensions orthogonal to the $(\delta + 3)$-brane:

$$r' \sim \frac{1}{M_s} \left[\alpha_s(M_s) \left(M_s^\delta R_1 ... R_\delta \right)^{1/2} \frac{M_{\rm Planck}}{M_s} \right]^{\frac{2}{(6-\delta)}} . \tag{14}$$

Using the value $\alpha_s(M_s) \sim 1/50$ for the gauge coupling at the string scale, $M_{\rm Planck} \sim 10^{19}$ GeV, and $R_1 ... R_\delta \sim (100 \text{ TeV})^\delta$ gives $r' \sim 10^{-7}$ cm for $\delta = 4$, and $r' \sim 10^{-11}$ cm for $\delta = 3$. There are of course general issues that remain to be solved in the context of D-branes. For example, how to construct stable non-supersymmetric brane configurations, or what brane configurations correspond to the three generations of chiral fermions [18].

In conclusion, the top condensation seesaw mechanism is a compelling scenario for electroweak symmetry breaking. It is remarkable that viable models involving this mechanism do not require an extension of the Standard Model gauge group, provided there are a few compact dimensions accessible to the gluons. If the future collider experiments will probe the composite Higgs sector, we will be able to test the existence of the KK modes of the gluons, and consequently the structure of the space-time.

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