Towards a Realistic Model of Higgsless Electroweak Symmetry Breaking

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(Received 18 August 2003; published 10 March 2004)

We present a 5D gauge theory in warped space based on a bulk $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge group where the gauge symmetry is broken by boundary conditions. The symmetry breaking pattern and the mass spectrum resemble that in the standard model (SM). To leading order in the warp factor the ρ parameter and the coupling of the Z (S parameter) are as in the SM, while corrections are expected at the level of a percent. From the anti–de Sitter (AdS) conformal field theory point of view the model presented here can be viewed as the AdS dual of a (walking) technicolorlike theory, in the sense that it is the presence of the IR brane itself that breaks electroweak symmetry, and not a localized Higgs on the IR brane (which should be interpreted as a composite Higgs model). This model predicts the lightest W, Z, and γ resonances to be at around 1.2 TeV, and no fundamental (or composite) Higgs particles.

DOI: 10.1103/PhysRevLett.92.101802 PACS numbers: 12.60.–i, 11.10.Kk, 11.15.Ex

The last unresolved mystery of the standard model (SM) of particle physics is the mechanism for electroweak symmetry breaking (EWSB). Within the SM it is assumed that a fundamental Higgs scalar is responsible for EWSB. This particle has not been observed yet, and its presence raises other fundamental issues like the hierarchy problem (that is how to avoid large quantum corrections to the mass of a light scalar). Nevertheless, the presence of such a Higgs scalar seems to be necessary; otherwise the scattering amplitudes of the longitudinal components of the massive W and Z bosons would blow up at scales of the order of 1 TeV, indicating new strongly interacting physics.

Recently, in collaboration with Murayama, we reexamined [1] the issue of longitudinal gauge boson scattering and found that there might be an alternative way to unitarize the gauge boson scattering amplitudes without a Higgs, if there is a tower of massive Kaluza-Klein (KK) gauge bosons present in these theories. In [1] we presented a toy model implementing this idea based on an $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge symmetry in an extra dimension where the gauge symmetry is broken by boundary conditions (BC's). There we found that the gauge boson spectrum somewhat resembles that in the SM; however the ρ parameter deviated from unity by as much as 10%, and the lowest KK excitations of the W and Z were too light for the model to be considered realistic.

In this paper we consider a similar model in a warped Randall-Sundrum (RS) [2] extra dimensional background. The motivation for considering this modification comes from the AdS/CFT correspondence [3]. The main problem with the flat-space model was the massive violation of custodial SU(2) symmetry which is manifested in the large deviation of ρ from one; therefore one would like to ensure that custodial SU(2) be maintained to leading order. A possible solution to this problem in the context of anti–de Sitter (AdS) space has been recently pointed out by Agashe *et al.* [4]. If one considers an AdS background, then one has a dual interpretation of the theory in terms of a spontaneously broken conformal field theory (CFT): the breaking of the conformal invariance is manifested by the presence of an infrared (TeV) brane, and the fields localized on the TeV brane are interpreted as bound states of the CFT. Gauge fields in the bulk correspond to global symmetries (that are weakly gauged) on the CFT side. This means that the $SU(2)_L \times SU(2)_R$ gauge symmetry in the bulk will ensure the presence of custodial SU(2) on the CFT side [4]. The symmetry breaking pattern on the TeV brane is $SU(2)_L \times SU(2)_R \to SU(2)_D$, which is exactly as in the SM, and preserves custodial isospin. The main difference between this model and other RS models with gauge fields in the bulk (such as [4,5]) is that here electroweak symmetry is broken by the presence of the TeV brane itself, rather than by a scalar Higgs localized on the TeV brane. The models with a TeV brane localized Higgs should be interpreted as the duals of composite Higgs models, where there is a scalar bound state of the strongly interacting CFT that is responsible for electroweak symmetry breaking. On the other hand, the model under consideration here, where electroweak symmetry breaking is due to the BC's on the TeV brane, should be interpreted as the dual of a (walking) technicolorlike theory, since it is the strong dynamics itself (the appearance of the TeV brane) that breaks the electroweak symmetry. Note that in the AdS picture one can interpolate between the technicolor and the composite Higgs models by dialing the expectation value of a brane localized Higgs field. For very large vacuum expectation values (VEV)'s the Higgs expels the wave functions and becomes a theory with BC breaking of electroweak symmetries corresponding to technicolor, while for small VEV's one gets the usual RS picture corresponding to a composite Higgs model.

The $SU(2)_R \times U(1)_{B-L}$ symmetry has to be broken in the UV to ensure that one has the right electroweak group at low energies. This can again be achieved by a BC breaking on the Planck brane, but will have the effect of giving corrections to electroweak observables. In the limit when the warp factor becomes infinitely large (the Planck brane is moved to the boundary of AdS) these corrections will vanish, but for a finite warp factor they will be suppressed by the log of the warp factor. These are relatively small compared to the flat-space model considered in [1], but are still about the order of the experimental precision of the electroweak observables. Therefore these corrections may still turn out to be too large, but this requires a detailed calculation of the electroweak precision observables including loop corrections from the relatively light KK excitations (and excluding the SM Higgs loops) to decide whether this particular model can be completely realistic or not. Either way, we consider the fact that the lowest order predictions reproduce the SM results without a Higgs to be a confirmation that the ideas presented in [1] could perhaps be implemented in a realistic way.

We want to study the possibility of breaking the electroweak symmetry $SU(2)_L \times U(1)_Y$ down to $U(1)_Q$ by BC's without relying on a Higgs mechanism in the bulk. We will consider a bulk $SO(4) \times U(1)_{B-L} \sim SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge group compactified in a warped RS background [2]. We will use the conformally flat metric

$$ds^2 = \left(\frac{R}{z}\right)^2 (\eta_{\mu\nu} dx^\mu dx^\nu - dz^2), \qquad (1)$$

where z is on the interval $[R, R']$. In RS-type models, R is typically $\sim 1/M_{\text{Pl}}$ and $R' \sim \text{TeV}^{-1}$. On the TeV brane at $z = R'$ we break SO(4) down to $SU(2)_D$. On the Planck brane, $z = R$, we break $SU(2)_R \times U(1)_{B-L}$ down to the usual $U(1)_Y$ hypercharge. Thus in the end only $U(1)_Q$, corresponding to electromagnetism, remains unbroken. We denote by $A_M^{R\,a}$, $A_M^{L\,a}$, and B_M the gauge bosons of $SU(2)_R$, $SU(2)_L$, and $U(1)_{B-L}$, respectively; g_5 is the gauge coupling of the two SU(2)'s and $\tilde{g}_5$, the gauge coupling of $U(1)_{B-L}$. We impose the following BC's $[A_M^{\pm\,a} = (A_M^{L\,a} \pm A_M^{R\,a})/\sqrt{2}]$:

$$\text{at} \quad z = R' : \begin{cases} \partial_z A_\mu^{+a} = 0, & A_\mu^{-a} = 0, \quad \partial_z B_\mu = 0, \\ A_5^{+a} = 0, & \partial_z A_5^{-a} = 0, \quad B_5 = 0, \end{cases} \qquad (2)$$

at

$$z = R : \begin{cases} \partial_5 A_\mu^{L\,a} = 0, & A_\mu^{R\,1,2} = 0, \\ \partial_z (g_5 B_\mu + \tilde{g}_5 A_\mu^{R\,3}) = 0, & \\ \tilde{g}_5 B_\mu - g_5 A_\mu^{R\,3} = 0, & \\ A_5^{L\,a} = 0, & A_5^{R\,a} = 0, \quad B_5 = 0. \end{cases} \qquad (3)$$

These BC's can be thought of as arising from Higgses on each brane in the limit of large VEVs which decouples the Higgs from gauge boson scattering [1]. The Higgs on the TeV brane is a bifundamental under the two SU(2)'s, while the Higgs on the Planck brane is a fundamental under $SU(2)_R$ and has charge $1/2$ under $U(1)_{B-L}$ so that a VEV in the lower component preserves $Y = T_3 + B - L$.

The KK mode expansion for the gauge fields in this background is given by

$$\psi_k^{(A)}(z) = z[a_k^{(A)} J_1(q_k z) + b_k^{(A)} Y_1(q_k z)], \qquad (4)$$

where A labels the corresponding gauge boson. Because of the mixing of the various gauge groups, the KK decomposition is slightly complicated but it is obtained by simply enforcing the BC's:

$$B_\mu = g_5 a_0 \gamma_\mu(x) + \sum_{k=1}^{\infty} \psi_k^{(B)}(z) Z_\mu^{(k)}(x), \qquad (5)$$

$$A_\mu^{L,R\,3} = \tilde{g}_5 a_0 \gamma_\mu(x) + \sum_{k=1}^{\infty} \psi_k^{(L,R\,3)}(z) Z_\mu^{(k)}(x), \qquad (6)$$

$$A_\mu^{L,R\,\pm} = \sum_{k=1}^{\infty} \psi_k^{(L,R\,\pm)}(z) W_\mu^{(k)\,\pm}(x). \qquad (7)$$

Here $\gamma(x)$ is the 4D photon, which has a flat wave function due to the unbroken $U(1)_Q$ symmetry, and $W_\mu^{(k)\,\pm}(x)$ and $Z_\mu^{(k)}(x)$ are the KK towers of the massive W and Z gauge bosons, the lowest of which are supposed to correspond to the observed W and Z.

The equation determining the tower of W masses can be read by substituting (7) into the BC's:

$$(R_0 - \tilde{R}_0)(R_1 - \tilde{R}_1) + (\tilde{R}_1 - R_0)(\tilde{R}_0 - R_1) = 0, \qquad (8)$$

where the ratios $R_{0,1}$ and $\tilde{R}_{0,1}$ are given by $R_i \equiv Y_i(MR)/J_i(MR)$, $\tilde{R}_i \equiv Y_i(MR')/J_i(MR')$. To leading order in $1/R$ and for $\log(R'/R) \gg 1$, the lightest solution for this equation for the mass of the $W^\pm$'s is

$$M_W^2 = \frac{1}{R'^2 \log(\frac{R'}{R})}. \qquad (9)$$

Note that this result does not depend on the 5D gauge coupling, but only on the scales R, R'. Taking $R = 10^{-19}\ \text{GeV}^{-1}$ will fix $R' = 2 \times 10^{-3}\ \text{GeV}^{-1}$.

The equation determining the masses of the KK tower for the Z (the states that are mostly A^{L3} or A^{R3}) is given by

$$g_5^2[(R_0 - \tilde{R}_0)(R_1 - \tilde{R}_1) + (\tilde{R}_1 - R_0)(\tilde{R}_0 - R_1)] = 2\tilde{g}_5^2 (R_0 - \tilde{R}_1)(\tilde{R}_0 - R_1). \qquad (10)$$

The lowest mass of the Z tower is approximately given by

$$M_Z^2 = \frac{g_5^2 + 2\tilde{g}_5^2}{g_5^2 + \tilde{g}_5^2} \frac{1}{R'^2 \log(\frac{R'}{R})}. \qquad (11)$$

Finally, on top of the massless photon, there is a third tower of states, corresponding to the excited modes of the photon (the particles that are mostly B type), whose masses are given by

$$R_0 = \tilde{R}_0. \tag{12}$$

In order to check whether these predictions agree with those of the SM we need to relate the bulk couplings g_5, $\tilde{g}_5$ to the effective SM couplings g, g'. This has to be done by introducing matter fields. Locally at the Planck brane ($z = R$ boundary), a $SU(2)_L \times U(1)_Y$ subgroup remains unbroken. We can introduce matter fields localized on this boundary. For simplicity consider first a scalar $SU(2)_L$ doublet with a $U(1)_{B-L}$ charge q. Its interactions with the gauge boson KK modes are generated through the Planck brane localized covariant derivative. Using the KK decomposition (7), one can evaluate this covariant derivative, which needs to be matched to the SM expression of the coupling of an $SU(2)_L$ doublet with hypercharge q. To be able to identify the first massive KK gauge bosons $Z^{(1)}$ and $W^{(1)}$ with the SM Z and W, we need to determine the gauge boson wave functions on the Planck brane and the integral of the square of the wave function in order to determine the normalization. To leading order (for $R \gg R'$) the integrals are dominated by the region near the Planck brane ($z \sim R$), so in fact the wave functions on the Planck brane are all that is needed. More specifically, from the expansion for small arguments of the Bessel functions appearing in (4), the wave function of a mode with mass $M \ll 1/R'$ can be written as [6]

$$\psi(z) \approx c_0 + M^2 z^2\left(c_1 - \frac{c_0}{2}\log\frac{z}{R}\right) + \mathcal{O}(M^4 z^4), \tag{13}$$

with c_0 at most of order 1, and c_1 at most of order $\mathcal{O}(\log(R'/R))$. Thus in the leading-log approximation

$$\int_R^{R'} dz\left(\frac{R}{z}\right)\psi(z)^2 \approx R\, c_0^2 \log\left(\frac{R'}{R}\right). \tag{14}$$

The boundary conditions on the bulk gauge fields give the following results for the leading term in the wave function for the lightest charged gauge bosons:

$$c_0^{(L\pm)} = c_\pm, \qquad c_0^{(R\pm)} \approx 0, \tag{15}$$

while for the neutral gauge bosons we find in the same approximation

$$c_0^{(L3)} \approx -c, \qquad c_0^{(R3)} \approx \frac{c\tilde{g}_5^2}{g_5^2 + \tilde{g}_5^2}, \qquad c_0^{(B)} \approx \frac{c\, g_5\, \tilde{g}_5}{g_5^2 + \tilde{g}_5^2}. \tag{16}$$

Using these results it can be checked that the usual SM relations are exactly satisfied (to leading-log order) and, from the coupling of the photon and the W, we can identify the 4D SM couplings in terms of the 5D gauge couplings by

$$g^2 = \frac{g_5^2}{R\log(R'/R)}, \tag{17}$$

$$e^2 = \frac{g_5^2 \tilde{g}_5^2}{(g_5^2 + 2\tilde{g}_5^2)R\log(R'/R)}. \tag{18}$$

One can also check that the full SM structures of the Z couplings are also reproduced to this approximation. Hence the ρ parameter in the leading-log approximation is

$$\rho = \frac{M_W^2}{M_Z^2 \cos^2\theta_W} \approx 1. \tag{19}$$

Note that the fact that the full structure of the SM coupling is reproduced implies that at the leading-log level there is no S parameter either. An S parameter in this language would have manifested itself in an overall shift of the coupling of the Z compared to its SM value evaluated from the W and γ couplings, which are absent at this order of approximation. The corrections to the SM relations will appear in the next order of the log expansion and are expected to be of the order of a percent. To evaluate the predictions of this model to a precision required by the measurements of the electroweak observables one needs to calculate at least the next order of corrections to the masses and couplings, together with the loop effects of the KK gauge bosons, and subtract the usual Higgs contributions.

The next issue: What are the masses of the KK excitations of the W and Z? One can see by numerically solving Eqs. (8) and (10) that $M_2^W \sim M_2^Z \sim M_2^\gamma \sim 1.2$ TeV. In terms of an energy expansion, the E^4 terms of the longitudinal WW scattering would blow up at energies of a few hundred GeV in the absence of a Higgs doublet; however to cancel those the effective four-point vertex obtained from integrating out a heavy W' and Z' is sufficient. The E^2 amplitudes would blow up at 1.8 TeV, which can be unitarized by the appearance of these new states. The next set of resonances arise at $M_3^W \sim M_3^Z \sim 1.9$ TeV. These masses are high enough to have evaded direct detection at the Tevatron, but should be within the reach of the Large Hadron Collider.

In the SM the Higgs is used not only to break electroweak symmetry, but also to generate fermion masses. For technicolor theories this generically poses a serious problem. In this model, the fermions can be added as bulk fermions that are doublets of $SU(2)_L$ (the left-handed fermions) and of $SU(2)_R$ (the right-handed fermions). Bulk fermions are generically Dirac fermions, however on an interval in warped space only one of the chiralities will have a zero mode. The location of the zero mode in warped space depends on the bulk mass term [7] and can be localized close to the Planck brane for the first two generations and the third generation leptons, which will imply that the gauge couplings for these fields will be as assumed above. For the right-handed top quark, one can localize the wave function of the zero mode closer to the TeV brane.

Since the theory on the TeV brane is vectorlike, a mass for the zero modes can be added on the TeV brane, which corresponds to a dynamical isospin symmetric fermion mass in the CFT language. The size of the physical mass will then depend on the location of the zero mode and the value of the mass term on the TeV brane. However because of the unbroken $SU(2)_D$ symmetry on the TeV brane these masses must be isospin symmetric; that is the masses for the up- and down-type quarks are equal at this point. Isospin splitting can be introduced for the leptons via Majorana masses on the Planck brane for the right-handed neutrinos using the seesaw mechanism, and via Dirac mass mixing with extra $SU(2)_R$ singlet fermions on the Planck brane [$SU(2)_R$ is broken on that brane so isospin breaking is allowed]. For the quarks this will effectively yield a top-quark seesaw-type model for the mass spectrum.

In summary, we have presented a 5D model in warped space where electroweak symmetry is broken by boundary conditions. The leading order predictions for the mass spectrum and coupling of the gauge bosons agree with the SM results, and the first excited W and Z fields appear at around a TeV, which is low enough to unitarize the scattering amplitudes. This model can be viewed as the AdS dual of a walking technicolorlike theory, and as such one needs to calculate the leading corrections to electroweak precision observables, which are estimated to be of the order of a percent.

We thank K. Agashe, S. Chivukula, A. Cohen, J. Hubisz, A. Katz, M. Luty, R. Rattazzi, Y. Shirman, L. Simmons, and R. Sundrum for useful discussions and comments. We thank the Aspen Center for Physics for its hospitality to us while this work was in progress. C. C. also thanks the T-8 group at Los Alamos for their hospitality while working on this project. C. C. and C. G. also thank the KITP at UC Santa Barbara for their hospitality while finishing this project. The research done at the KITP is supported by the NSF under Grant No. PHY99-07949.

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