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Towards a Realistic Model of Higgsless Electroweak Symmetry Breaking

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Abstract

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 resembles that in the standard model (SM). To leading order in the warp factor the ρ parameter and the coupling of the Z (or equivalently the S-parameter) are as in the SM, while corrections are expected at the level of a percent. From the AdS/CFT point of view the model presented here can be viewed as the AdS dual of a (walking) technicolor-like 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.

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 order 1 TeV, indicating new strongly interacting physics.

Recently, in collaboration with H. Murayama, we have re-examined [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 (see also [2, 3] and, for similar considerations in gravitational theories, see [4]). In [1] we have 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 ten percent, 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) [5] extra dimensional background. The motivation for considering this modification comes from the AdS/CFT correspondence [6]. 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. [7]. 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 [7]. 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 [7, 8, 9], see also [10]) 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) technicolor-like theory [11], 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 composite Higgs models by dialing the expectation value of a brane localized Higgs field.

For very large 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 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.* The BC breaking is equivalent to a Higgs mechanism on the brane in the limit of very large VEV's for the brane Higgs fields. These large VEV's will repel the wave functions of massive modes and ensure that the Higgs decouples from the gauge and matter fields [1]. As in [1] 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 (where $U(1)_{B-L}$ corresponds to gauging baryon minus lepton number), except here we will consider the theory compactified in a warped RS background [5]. We will use the conformally flat metric

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

where z is on the interval $[R, R']$. In RS-type models R is typically $\sim 1/M_{Pl}$ and $R' \sim \text{TeV}^{-1}$. On the TeV brane at $z = R'$ we break $SO(4)$ down to $SU(2)_D$ by Neumann and Dirichlet BC's. As discussed above, this corresponds, from the AdS/CFT point of view, to breaking the $SU(2)_L \times SU(2)_R$ global symmetry of the CFT to the diagonal subgroup, just as would happen in technicolor models. On the Planck brane, $z = R$, we break $SU(2)_R \times U(1)_{B-L}$ down to the usual $U(1)_Y$ hypercharge again by Neumann and Dirichlet BC's, to ensure that the low-energy gauge group without electroweak symmetry breaking is $SU(2)_L \times U(1)_Y$. 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

*Other interesting possibilities for EWSB using extra dimensions is to have the Higgs be the extra dimensional component of a gauge field, see for example [12], or to have a warped compactification where the would-be zero mode for the gauge field is not normalizable [13].

the following BC's:

$$\text{at } z = R': \quad \begin{cases} \partial_z (A^{L\,a}_\mu + A^{R\,a}_\mu) = 0,\ A^{L\,a}_\mu - A^{R\,a}_\mu = 0,\ \partial_z B_\mu = 0, \\ (A^{L\,a}_5 + A^{R\,a}_5) = 0,\ \partial_z (A^{L\,a}_5 - A^{R\,a}_5) = 0,\ B_5 = 0. \end{cases} \tag{2}$$

$$\text{at } z = R: \quad \begin{cases} \partial_5 A^{L\,a}_\mu = 0,\ A^{R\,1,2}_\mu = 0, \\ \partial_z (g_5 B_\mu + \tilde{g}_5 A^{R\,3}_\mu) = 0,\ \tilde{g}_5 B_\mu - g_5 A^{R\,3}_\mu = 0, \\ A^{L\,a}_5 = 0,\ A^{R\,a}_5 = 0,\ B_5 = 0. \end{cases} \tag{3}$$

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 bi-fundamental 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 $Y = T_3 + B - L$. If one insists on the breaking by BC picture, and does not want to think of it as Higgs fields on the branes with large VEV's one could ask why the gauge couplings would ever appear in the BC's. However, that is clearly dependent on what normalization is chosen for the gauge fields. Above we have chosen 5D canonical normalization. If we went to the "more natural" normalization where the gauge couplings appear only in front of the gauge kinetic term in the denominator, then the BC's on the Planck brane $z = R$ would simply be $\partial_z (B_\mu + A^{R\,3}_\mu) = 0$, and $B_\mu = A^{R\,3}_\mu$. Thus it does not seem unnatural for the gauge couplings to appear in the expressions for the wave functions once we go back to a canonically normalized basis.

We will work in unitary gauge (the $\xi \to \infty$ limit of the R_ξ gauge discussed in detail in [1]). Since the 5th components of the gauge fields do not have zero modes, they will all decouple in the unitary gauge. The Euclidean bulk equation of motion satisfied by spin-1 fields in AdS space is

$$(\partial_z^2 - \frac{1}{z}\partial_z + q^2)\psi(z) = 0, \tag{4}$$

where the solutions in the bulk are assumed to be of the form $A_\mu(q)e^{-iqx}\psi(z)$. The the KK mode expansion is given by the solutions to this equation which are of the form

$$\psi_k^{(A)}(z) = z\left(a_k^{(A)} J_1(q_k z) + b_k^{(A)} Y_1(q_k z)\right)\ , \tag{5}$$

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

$$
\begin{aligned}
B_\mu(x,z) &= g_5\, a_0 \gamma_\mu(x) + \sum_{k=1}^{\infty} \psi_k^{(B)}(z)\, Z_\mu^{(k)}(x)\,, && (6)\\
A_\mu^{L\,3}(x,z) &= \tilde g_5\, a_0 \gamma_\mu(x) + \sum_{k=1}^{\infty} \psi_k^{(L3)}(z)\, Z_\mu^{(k)}(x)\,, && (7)\\
A_\mu^{R\,3}(x,z) &= \tilde g_5\, a_0 \gamma_\mu(x) + \sum_{k=1}^{\infty} \psi_k^{(R3)}(z)\, Z_\mu^{(k)}(x)\,, && (8)\\
A_\mu^{L\,\pm}(x,z) &= \sum_{k=1}^{\infty} \psi_k^{(L\pm)}(z)\, W_\mu^{(k)\,\pm}(x)\,, && (9)\\
A_\mu^{R\,\pm}(x,z) &= \sum_{k=1}^{\infty} \psi_k^{(R\pm)}(z)\, W_\mu^{(k)\,\pm}(x)\,. && (10)
\end{aligned}
$$

Here $\gamma(x)$ is the 4D photon, which has a flat wavefunction 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 of by substituting (9-10) into the BC's, and is given by

$$(R_0 - \tilde R_0)(R_1 - \tilde R_1) + (\tilde R_1 - R_0)(\tilde R_0 - R_1) = 0 \tag{11}$$

where the ratios $R_{0,1}$ and $\tilde R_{0,1}$ are given by

$$R_i \equiv \frac{Y_i(MR)}{J_i(MR)},\;\; \tilde R_i \equiv \frac{Y_i(MR')}{J_i(MR')}. \tag{12}$$

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

$$M_W^2 = \frac{1}{R'^2 \log\left(\frac{R'}{R}\right)}\,. \tag{13}$$

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

$$2\tilde g_5^2 (R_0 - \tilde R_1)(\tilde R_0 - R_1) = g_5^2 \left[(R_0 - \tilde R_0)(R_1 - \tilde R_1) + (\tilde R_1 - R_0)(\tilde R_0 - R_1)\right] \tag{14}$$

The lowest mass of the Z tower is approximately given by

$$M_Z^2 = \frac{g_5^2 + 2\tilde g_5^2}{g_5^2 + \tilde g_5^2} \frac{1}{R'^2 \log\left(\frac{R'}{R}\right)}\,. \tag{15}$$

Finally, 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

$$R_0 = \tilde{R}_0. \tag{16}$$

This does not have a light mode (the zero mode corresponding to the massless photon has been separated out explicitly in (6-8)).

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 ($z = R$), 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 localized covariant derivative

$$D_\mu \Phi = \partial_\mu \Phi - \frac{i}{2} \begin{pmatrix} 2\tilde{g}_5 q B_\mu + g_5 A_\mu^{L\,3} & g_5 (A_\mu^{L\,1} - i A_\mu^{L\,2}) \\ g_5 (A_\mu^{L\,1} + i A_\mu^{L\,2}) & 2\tilde{g}_5 q B_\mu - g_5 A_\mu^{L\,3} \end{pmatrix}_{|z=R} \Phi. \tag{17}$$

Using the KK decomposition (6)-(10), we evaluate the gauge fields at the boundary and the scalar covariant derivative becomes

$$\begin{aligned} D_\mu \Phi = \partial_\mu \Phi \; & - \; i g_5 \tilde{g}_5 a_0 \begin{pmatrix} q + \frac{1}{2} & 0 \\ 0 & q - \frac{1}{2} \end{pmatrix} \gamma_\mu \Phi \\ & - \; i \sum_{k=1}^{\infty} \begin{pmatrix} \tilde{g}_5 q \psi_k^{(B)} + \frac{1}{2} g_5 \psi_k^{(L3)} & 0 \\ 0 & \tilde{g}_5 q \psi_k^{(B)} - \frac{1}{2} g_5 \psi_k^{(L3)} \end{pmatrix}_{|z=R} Z_\mu^{(k)} \Phi \\ & - \; \sum_{k=1}^{\infty} i g_5 \tfrac{1}{\sqrt{2}}\, \psi_{k\,|z=R}^{(L\pm)} \begin{pmatrix} 0 & W_\mu^{(k)\,+} \\ W_\mu^{(k)\,-} & 0 \end{pmatrix} \Phi. \end{aligned} \tag{18}$$

This needs to be matched to the SM expression of the coupling of an $SU(2)_L$ doublet with hypercharge q which is given by

$$D_\mu \Phi = \partial_\mu \Phi - i \begin{pmatrix} \frac{g^2 - 2qg'^2}{2\sqrt{g^2+g'^2}} Z_\mu + e(q + \frac{1}{2}) \gamma_\mu & \frac{g}{\sqrt{2}}\, W_\mu^+ \\ \frac{g}{\sqrt{2}}\, W_\mu^- & \frac{-g^2 - 2qg'^2}{2\sqrt{g^2+g'^2}} Z_\mu + e(q - \frac{1}{2}) \gamma_\mu \end{pmatrix} \Phi. \tag{19}$$

To be able to identify the first massive KK gauge bosons $Z^{(1)}$ and $W^{(1)}$ with the SM Z and W we then need to determine the gauge boson wavefunctions on the Planck brane and the integral of the square of the wavefunction 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 wavefunctions on the Planck brane are all that is needed. More specifically, from the expansion for small arguments of the Bessel functions appearing in (5), the wavefunction of a mode with mass $M \ll 1/R'$ can be written as [9]:

$$\psi(z) \approx c_0 + M^2 \left(c_1 z^2 - \frac{c_0}{2} z^2 \log(z/R) \right) + \mathcal{O}(M^4 z^4), \tag{20}$$

with c_0 at most of order one, and c_1 at most of order $\mathcal{O}(\log(R'/R))$. Thus

$$\int_R^{R'} dz \left(\frac{R}{z}\right) \psi(z)^2 = R\left[c_0^2 \log\left(\frac{R'}{R}\right) + M^2 c_0 c_1 R'^2 - \frac{1}{4} M^2 c_0^2 R'^2 (2\log\frac{R'}{R} - 1) + ...\right]. \quad (21)$$

As we have seen above, $M^2 = \mathcal{O}(1/\log(R'/R))$, so in the leading-log approximation:

$$\int_R^{R'} dz \left(\frac{R}{z}\right) \psi(z)^2 \approx R\, c_0^2 \log\left(\frac{R'}{R}\right) \ . \quad (22)$$

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

$$c_0^{(L\pm)} = c_\pm \ , \quad c_0^{(R\pm)} \approx 0 \ , \quad (23)$$

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

$$c_0^{(L3)} \approx -\, c \ , \quad c_0^{(R3)} \approx c\, \frac{\tilde{g}_5^2}{g_5^2 + \tilde{g}_5^2} \ , \quad c_0^{(B)} \approx c\, \frac{g_5\, \tilde{g}_5}{g_5^2 + \tilde{g}_5^2} \ . \quad (24)$$

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

$$g^2 = \frac{g_5^2 \psi^{(L\pm)}(R)^2}{\int_R^{R'} dz \left(\frac{R}{z}\right) (\psi^{(L\pm)}(z)^2 + \psi^{(R\pm)}(z)^2)} = \frac{g_5^2}{R\log(R'/R)}, \quad (25)$$

$$e^2 = \frac{\tilde{g}_5^2 g_5^2 a_0^2}{\int_R^{R'} dz \left(\frac{R}{z}\right) (2\tilde{g}_5^2 + g_5^2) a_0^2} = \frac{g_5^2 \tilde{g}_5^2}{(g_5^2 + 2\tilde{g}_5^2) R\log(R'/R)}. \quad (26)$$

The full SM structure of the Z couplings are also reproduced since:

$$g^2 \cos\theta_W^2 = \frac{g_5^2 \psi^{(L3)}(R)^2}{\int_R^{R'} dz \left(\frac{R}{z}\right) (\psi^{(L3)2} + \psi^{(R3)2} + \psi^{(B)2})} = \frac{g_5^2}{R\log(R'/R)} \frac{g_5^2 + \tilde{g}_5^2}{g_5^2 + 2\tilde{g}_5^2}, \quad (27)$$

$$g'^2 = \frac{g_5^2 \tilde{g}_5^2}{(g_5^2 + \tilde{g}_5^2) R\log(R'/R)}, \quad (28)$$

$$\sin\theta_W = \frac{\tilde{g}_5}{\sqrt{g_5^2 + 2\tilde{g}_5^2}} = \frac{g'}{\sqrt{g^2 + g'^2}}. \quad (29)$$

Hence the ρ parameter in the leading log approximation is

$$\rho = \frac{M_W^2}{M_Z^2 \cos^2\theta_W} \approx 1 \ . \quad (30)$$

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 according to (27)–(29) 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 is: what are the masses of the KK excitations of the W and Z? One can see by numerically solving Eqs. (12) and (14) (and using $R = 10^{-19}$ GeV^{-1}, $R' = 2 \cdot 10^{-3}$GeV^{-1}) that $M_2^W \sim M_2^Z \sim M_2^\gamma \sim 1.2$ TeV. These masses are high enough to have evaded direct detection at the Tevatron, but should be within the reach of the LHC. In terms of an energy expansion, the E^4 terms of the longitudinal WW scattering would blow up at energies of 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.

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 of the SM). Bulk fermions are generically Dirac fermions, however on an interval in warped space only one of the chiralities will have a zero mode [14]. The location of the zero mode in warped space depends on the bulk mass term [15], 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 vector-like (only $SU(2)_D$ is unbroken there), 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 mass for the up and down type quarks (and similarly for the charged leptons and neutrinos) 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 see-saw mechanism, and via Dirac mass mixing with extra $SU(2)_R$ singlet fermions on the Planck brane (isospin breaking can be introduced there since $SU(2)_R$ is broken on that brane). For the quarks this will effectively yield a top-quark see-saw 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 technicolor-like theory, and as such one needs to calculate the leading corrections to electroweak precision observables, which are estimated to be of order of a percent.

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