FERMILAB-Pub-03/284-T

Higgsless Theory of Electroweak Symmetry Breaking from Warped Space

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Abstract

We study a theory of electroweak symmetry breaking without a Higgs boson, recently suggested by Csaki *et al.* The theory is formulated in 5D warped space with the gauge bosons and matter fields propagating in the bulk. In the 4D dual picture, the theory appears as the standard model without a Higgs field, but with an extra gauge group G which becomes strong at the TeV scale. The strong dynamics of G breaks the electroweak symmetry, giving the masses for the W and Z bosons and the quarks and leptons. We study corrections in 5D which are logarithmically enhanced by the large mass ratio between the Planck and weak scales, and show that they do not destroy the structure of the electroweak gauge sector at the leading order. We introduce a new parameter, the ratio between the two bulk gauge couplings, into the theory and find that it allows us to control the scale of new physics. We also present a potentially realistic theory accommodating quarks and leptons and discuss its implications, including the violation of universality in the W and Z boson couplings to matter and the spectrum of the Kaluza-Klein excitations of the gauge bosons. The theory reproduces many successful features of the standard model, although some cancellations may still be needed to satisfy constraints from the precision electroweak data.

1 Introduction

Despite its extraordinary successes in describing physics below the energy scale of ~ 100 GeV, the standard model of particle physics does not address the mechanism of electroweak symmetry breaking. In the standard model the electroweak symmetry is broken by a vacuum expectation value (VEV) of a Higgs field, which is driven by a non-trivial potential introduced just to trigger the symmetry breaking. Moreover, this potential is not stable against radiative corrections and is extremely sensitive to any high energy scales, leading to the notorious gauge hierarchy problem. An elegant idea to address this problem is to assume that the electroweak symmetry is broken by some gauge dynamics whose interaction becomes strong at the TeV scale [1]. This idea has been pursued over 25 years, but the attempts of constructing realistic theories have encountered many obstacles such as the difficulties of obtaining realistic quark and lepton masses, suppressing unwanted flavor violation, and complying with precision electroweak data.

While the problems listed above depend in principle on the dynamics of the strong gauge sector and there may be some theories free from these problems, the strong coupling dynamics makes it difficult to address these issues quantitatively and construct realistic theories in which the agreement with experiments is reliably seen. This is particularly problematic because the naive estimate based on an analogy with QCD suggests that generic theories of dynamical electroweak symmetry breaking are in disagreement with data [2]. There are various attempts for overcoming these obstacles, but none seems completely satisfactory.

The AdS/CFT duality [3] offers a new possibility of addressing these questions. Suppose the gauge interaction responsible for electroweak symmetry breaking has a large number of "colors" and is almost conformal at energies larger than the dynamical scale. In this case the theory has an equivalent description in terms of a weakly coupled 5D theory defined on the truncated anti de-Sitter (AdS) space [4]. The geometry of the 5D theory is that of Randall and Sundrum [5], and in this picture the hierarchy between the Planck and the electroweak scales is understood in terms of the AdS warp factor. The original picture of dynamical electroweak symmetry breaking is then mapped to the following situation in 5D. The electroweak gauge symmetry, with the gauge fields propagating in the bulk, is broken at the infrared brane by either boundary conditions or a VEV of the order of the local cutoff scale. This line of constructing theories of dynamical electroweak symmetry breaking has recently been considered in [6] with a new ingredient of introducing a gauged custodial symmetry in the bulk of a 5D theory [7], which corresponds in the 4D picture to imposing a global custodial symmetry on the strong dynamics.[1] It has been claimed in [6] that the correct masses and couplings of the electroweak

[1]Models of a composite Higgs boson using AdS/CFT have been considered in [8, 9, 7].

gauge bosons are reproduced at the level of a percent. In this paper we first study this issue, including effects in 5D which are logarithmically enhanced by the large mass ratio between the Planck and the weak scales. We explicitly show that these corrections do not destroy the structure of the electroweak gauge sector obtained in [6] at the leading order.

We also extend the analysis of [6] to the case where the couplings of the bulk gauge groups take arbitrary values. This allows us to incorporate a realistic structure for quarks and leptons relatively easily and to make the scale of new physics higher. While we do not perform a detailed analysis for the corrections to the electroweak observables at a level of a percent, we find that this extra freedom for pushing up the scale of new physics allows us to reduce the deviations from the standard model. In view of these, we think it is quite interesting if we could construct a theory of electroweak symmetry breaking that does not rely on the Higgs mechanism but reproduces successful features of the standard model, including the fermion mass generation and smallness of flavor changing neutral currents (FCNCs). We thus attempt to construct such theories.

We formulate the theory in the 5D AdS space with the gauge group $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_X$ in the bulk. These gauge groups are broken to $SU(3)_C \times U(1)_{\rm EM}$ by large VEVs for scalar fields located on the branes. Although we formally use scalar field VEVs to break gauge groups, the physical Higgs bosons decouple when these VEVs are large so that the theory does not contain any narrow scalar resonance. To incorporate the observed structure for the quark and lepton masses in a simple way, we choose representations for quarks and leptons asymmetric under the interchange between $SU(2)_L$ and $SU(2)_R$. We find that the theory does not give an immediate contradiction with existing experimental data for certain parameter region, although some amount of tuning may be necessary to satisfy precision electroweak constraints. We also consider the Kaluza-Klein (KK) towers for the gauge bosons and find that they are within the reach of the LHC.

The organization of the paper is as follows. In the next section we summarize the theory and framework. In section 3 we calculate the masses and couplings of the electroweak gauge bosons at the leading order in the warp factor and show that the structure of [6] is preserved despite the presence of corrections enhanced by a logarithm of the large mass ratio between the Planck and TeV scales. We give a generalized relation between the scale of new physics and the masses of W and Z, which allows us to push up the masses of the resonances in the sector of dynamical electroweak symmetry breaking. In section 4 we present a concrete model in the present context. We discuss fermion mass generation and FCNCs arising from the non-universality of fermion couplings to the electroweak gauge bosons. The spectrum for the KK towers of the gauge bosons are also discussed. Conclusions are given in section 5.

2 Theory and Framework

Following Refs. [7, 6] we consider an $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_X$ gauge theory in 5D warped space. The extra dimension is compactified on S^1/Z_2 and the metric is given by

$$ds^2 \equiv G_{MN} dx^M dx^N = e^{-2k|y|} \eta_{\mu\nu} dx^\mu dx^\nu + dy^2, \tag{1}$$

where y is the coordinate of the fifth dimension and the physical space is taken to be $0 \leq y \leq \pi R$; k is the AdS curvature scale. We take k to be around the 4D Planck scale, $k \sim M_{\rm pl}$, and $kR \sim 10$ to address the gauge hierarchy problem [5]. The gauge kinetic terms in the bulk are given by

$$\begin{aligned} S \;=\; \int d^4x \int dy \sqrt{-G} \Big[& -\frac{1}{4g_L^2} g^{MP} g^{NQ} \sum_{a=1}^{3} F_{MN}^{La} F_{PQ}^{La} - \frac{1}{4g_R^2} g^{MP} g^{NQ} \sum_{a=1}^{3} F_{MN}^{Ra} F_{PQ}^{Ra} \\ & -\frac{1}{4g_X^2} g^{MP} g^{NQ} F_{MN}^{X} F_{PQ}^{X} \Big], \end{aligned} \tag{2}$$

where F_{MN}^{La}, F_{MN}^{Ra} and F_{MN}^{X} are the field-strength tensors for $SU(2)_L$, $SU(2)_R$ and $U(1)_X$, and g_L, g_R and g_X are the 5D gauge couplings having mass dimensions $-1/2$; a is the indices for the adjoint representation of $SU(2)$. Here, we have suppressed $SU(3)_C$, since it is irrelevant for our discussion below.

To reduce the gauge symmetry down to $SU(3)_C \times U(1)_{\rm EM}$ at low energies, where $U(1)_{\rm EM}$ is the electromagnetism, we impose non-trivial boundary conditions on the gauge fields at $y = 0$ (the Planck brane) and $y = \pi R$ (the TeV brane). They are given by [6, 10]

$$\partial_y A_\mu^{La} = 0, \quad A_\mu^{R1,2} = 0, \quad \partial_y \Big(\frac{1}{g_R^2} A_\mu^{R3} + \frac{1}{g_X^2} A_\mu^{X} \Big) = 0, \quad A_\mu^{R3} - A_\mu^{X} = 0, \tag{3}$$

at $y = 0$ and

$$\partial_y \Big(\frac{1}{g_L^2} A_\mu^{La} + \frac{1}{g_R^2} A_\mu^{Ra} \Big) = 0, \quad A_\mu^{La} - A_\mu^{Ra} = 0, \quad \partial_y A_\mu^{X} = 0, \tag{4}$$

at $y = \pi R$; the A_5's obey Dirichlet (Neumann) boundary conditions if the corresponding A_μ's obey Neumann (Dirichlet) boundary conditions. Here, we have retained arbitrary bulk gauge couplings g_L, g_R and g_X for later purposes. In the 4D dual picture, this leads to the following situation. The theory below $k \sim M_{\rm pl}$ contains a gauge interaction with the group G, whose coupling evolves very slowly over a wide energy interval below k. This G gauge sector possesses a global $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_X$ symmetry whose $SU(3)_C \times SU(2)_L \times U(1)_Y$ subgroup is weakly gauged, where $U(1)_Y$ is a linear combination of $U(1)_X$ and the T_3 direction of $SU(2)_R$. Therefore, the theory in this energy interval appears as $SU(3)_C \times SU(2)_L \times U(1)_Y \times G$ gauge theory (with quarks and leptons transforming under $SU(3)_C \times SU(2)_L \times U(1)_Y$; see below). At the TeV scale the gauge interaction of G becomes strong. This dynamically breaks $SU(2)_L \times U(1)_Y$ down to $U(1)_{\rm EM}$, giving masses to the W and Z bosons (and to quarks and leptons).

The boundary conditions of Eqs. (3, 4) can also be viewed as the limiting case of the following brane Higgs breaking. Suppose that we have a Higgs field $\Sigma(\mathbf{1}, \mathbf{1}, \mathbf{2}, 1/2)$ on the Planck brane and $H(\mathbf{1}, \mathbf{2}, \mathbf{2}^*, 0)$ on the TeV brane, where the numbers in the parentheses represent gauge quantum numbers under $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_X$, and that these fields have VEVs:

$$\langle \Sigma \rangle = \begin{pmatrix} 0 \\ v_\Sigma \end{pmatrix}, \qquad \langle H \rangle = \begin{pmatrix} v_H & 0 \\ 0 & v_H \end{pmatrix}. \tag{5}$$

For large v_Σ and v_H (more precisely $v_\Sigma, v_H \gg (k/g_{L,R,X}^2)^{1/2}$) this Higgs breaking resembles the boundary condition breaking described above, because a large brane mass term for a gauge field pushes off the wavefunction of the gauge field from the brane and thus is equivalent to imposing the Dirichlet boundary condition [11]. The phenomenology becomes identical to that of the boundary condition breaking in the limit $v_\Sigma, v_H \to \infty$ as far as the gauge field sector is concerned. Note that the physical Higgs bosons arising from Σ and H decouple in this limit and no scalar particle remains in the spectrum. Our analyses below are applicable to both cases for the boundary condition breaking and the Higgs breaking described here.

To see whether the theory can be realistic or not, we have to consider the couplings of the gauge bosons to quarks and leptons. Here we consider the case of the Higgs breaking but the results should also apply to any potentially realistic theories with the boundary condition breaking. A single generation of quarks and leptons corresponds to:

$$\begin{array}{lll} q(\mathbf{3}, \mathbf{2}, \mathbf{1}, 1/6), & \bar{u}(\mathbf{3}^*, \mathbf{1}, \mathbf{2}, -1/6)|_{T_3^R=-1/2}, & \bar{d}(\mathbf{3}^*, \mathbf{1}, \mathbf{2}, -1/6)|_{T_3^R=1/2}, \\ l(\mathbf{1}, \mathbf{2}, \mathbf{1}, -1/2), & \bar{e}(\mathbf{1}, \mathbf{1}, \mathbf{2}, 1/2)|_{T_3^R=1/2}, & [\bar{n}(\mathbf{1}, \mathbf{1}, \mathbf{2}, 1/2)|_{T_3^R=-1/2}], \end{array} \tag{6}$$

where $q, \bar{u}, \bar{d}, l, \bar{e}$ and $\bar{n}$ are left-handed fermions and the numbers in the parentheses represent gauge quantum numbers under $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_X$; $T_3^R = \pm 1/2$ represents the $T_3 = \pm 1/2$ component of the $SU(2)_R$ doublet. A realistic structure for the quark and lepton sector is obtained, for example, as follows. We introduce the quark and lepton fields in the bulk, transforming as $q(\mathbf{3}, \mathbf{2}, \mathbf{1}, 1/6)$, $\psi_{\bar{u}}(\mathbf{3}^*, \mathbf{1}, \mathbf{2}, -1/6)$, $\psi_{\bar{d}}(\mathbf{3}^*, \mathbf{1}, \mathbf{2}, -1/6)$, $l(\mathbf{1}, \mathbf{2}, \mathbf{1}, -1/2)$ and $\psi_{\bar{e}}(\mathbf{1}, \mathbf{1}, \mathbf{2}, 1/2)$ [here we do not consider neutrino masses for simplicity]. They are Dirac fermions but due to the orbifold compactification on S^1/Z_2, only the left-handed fermions with the gauge quantum numbers given in the parentheses remain massless. The wavefunction profiles for these massless modes are controlled by the bulk mass terms for the fermions [12, 13], which we parameterize as $\mathcal{L}_{5\rm D} \supset -ck\bar{\Psi}\Psi$ where Ψ represents generic 5D (Dirac) fermions. For $c > 1/2$ ($c < 1/2$) the wavefunction for the left-handed zero-mode fermion is localized to the Planck

(TeV) brane. We take parameters c to be larger than $1/2$, at least for the first-two generation fermions. The unwanted zero modes, the $T_3^R = 1/2$ component of $\psi_{\bar{u}}$ and the $T_3^R = -1/2$ components of $\psi_{\bar{d}}$ and $\psi_{\bar{e}}$, are made heavy to get masses of order k, by introducing the Planck-brane localized left-handed fermions $\psi'_{\bar{u}}(\mathbf{3},\mathbf{1},\mathbf{1},-1/3)$, $\psi'_{\bar{d}}(\mathbf{3},\mathbf{1},\mathbf{1},2/3)$ and $\psi'_{\bar{e}}(\mathbf{1},\mathbf{1},\mathbf{1},0)$ and couplings $\delta(y)[\psi_{\bar{u}}\psi'_{\bar{u}}\Sigma + \psi_{\bar{d}}\psi'_{\bar{d}}\Sigma^\dagger + \psi_{\bar{e}}\psi'_{\bar{e}}\Sigma^\dagger]$. The low-energy matter content is then precisely that of the standard model quarks and leptons.

The bulk mass terms of $c > 1/2$ are required/beneficial for several reasons. First, since the electroweak gauge symmetry is broken by boundary conditions or a large VEV of the brane-localized Higgs, the wavefunctions of the W and Z bosons are not flat in the extra dimension. This generically introduces non-universality of the electroweak gauge couplings depending on the bulk fermion mass parameters, because the 4D gauge couplings are obtained by convolving the wavefunctions of the corresponding fermion with the gauge boson, which are not universal for fermions having different values of the bulk masses. This non-universality, however, is very small when c's are larger than $1/2$ (more detailed discussions are given in section 4). Thus we take $c > 1/2$ at least for the first-two generation fermions [c will be smaller for the right-handed top to give a large enough top quark mass]. Second, since fermion masses arise from operators localized on the TeV brane [bare mass terms in the case of boundary condition breaking, and masses through operators like $q\psi_{\bar{u}}H$, $q\psi_{\bar{d}}H$ and $l\psi_{\bar{e}}H$ in the case of Higgs breaking], having $c > 1/2$ suppresses the masses of the corresponding fermions, giving an understanding of the lightness of the first-two generation fermions through the localization of the fermion wavefunctions [13, 14]. Finally, by localizing the wavefunctions to the Planck brane, *i.e.* making c's larger than $1/2$, we can suppress effects from unwanted tree-level flavor-violating operators on the TeV brane. In any event, with c larger than $1/2$, the couplings of the quarks and leptons to the electroweak gauge bosons can be effectively read off from

$$\mathcal{D}_\mu \psi_L = \left[\partial_\mu \psi_L + \frac{i}{2} \begin{pmatrix} A_\mu^{L3} + 2qA_\mu^X & A_\mu^{L1} - iA_\mu^{L2} \\ A_\mu^{L1} + iA_\mu^{L2} & -A_\mu^{L3} + 2qA_\mu^X \end{pmatrix} \psi_L \right] \Bigg|_{y=0} \tag{7}$$

for left-handed fermions ψ_L transforming $(\mathbf{2},\mathbf{1},q)$ under $SU(2)_L \times SU(2)_R \times U(1)_X$, and from

$$\mathcal{D}_\mu \psi_R = \left[\partial_\mu \psi_R + \frac{i}{2} \begin{pmatrix} A_\mu^{R3} + 2qA_\mu^X & A_\mu^{R1} - iA_\mu^{R2} \\ A_\mu^{R1} + iA_\mu^{R2} & -A_\mu^{R3} + 2qA_\mu^X \end{pmatrix} \psi_R \right] \Bigg|_{y=0} \tag{8}$$

for right-handed fermions ψ_R transforming $(\mathbf{1},\mathbf{2},q)$ under $SU(2)_L \times SU(2)_R \times U(1)_X$.

It has been shown in [6] that, for fermion couplings of Eqs. (7, 8) and the gauge kinetic terms given by Eq. (2) with $g_L = g_R$, the tree-level structure of the standard model gauge couplings are reproduced. The relations among various γ, W and Z couplings and the masses of the W and Z bosons are then exactly those of the tree-level standard model, up to corrections of order $1/\pi kR = O(1\%)$. Although this level of agreement is not sufficient to satisfy constraints from precision electroweak data, we think it is interesting if a theory without a Higgs boson reproduced the standard model relations at this level. Thus we take the attitude that the theory is successful if it reproduced the standard model relations at the leading order in $1/\pi kR$. In the next section we extend the analysis of [6] to the case with logarithmically enhanced corrections in 5D and arbitrary bulk gauge couplings. We explicitly show that the successful structure in the electroweak sector is preserved under these extensions, and derive a generalized relation between the scale of new physics and the masses of W and Z.

3 Masses and Couplings of the Electroweak Gauge Bosons

Let us start by considering the 4D dual picture of the theory. In this picture, the gauge interaction G becomes strong at the TeV scale, producing resonances of masses of order TeV. These resonances have a tower structure. In particular, there are towers of spin-1 fields which have the quantum numbers of W, Z and γ. These towers then mix with the elementary gauge bosons of the weakly gauged $SU(2)_L$ and $U(1)_Y$ groups. The resulting spectrum consists of towers of gauge bosons with the quantum numbers of W and Z, whose lowest states are massive and identified as the standard-model W and Z bosons, and a tower of $U(1)$ gauge bosons, whose lowest mode is massless and identified with the photon. These towers of mass eigenstates are dual to the W, Z and γ KK towers in the 5D picture.

In the previous section, we have considered the theory in 5D with the gauge kinetic terms given by Eq. (2). This corresponds in the 4D picture to the following situation. In the absence of the G sector, the couplings of the elementary gauge bosons g_E are very large, $1/g_E^2 \simeq 0$. The observed values for the W, Z and γ couplings, $\alpha = g^2/4\pi$, are then obtained purely through the contribution from the G sector. In this case, the masses and couplings of W, Z and γ obtained after the elementary-composite mixings respect the tree-level standard-model relations at the leading order in $1/N$ and α, where N is the number of "colors" for G, since this corresponds in the 5D picture to deriving these quantities at tree level. The question is whether the above structure is preserved under corrections at higher orders in $1/N$ and α. Potentially dangerous ones are loops of elementary fields, as they induce non-zero values of $1/g_E^2$ and could change the structure. These corrections are loop suppressed but enhanced by a logarithm, so have a size of order $(c\alpha/4\pi)\ln(k/T)$, where c is a group theoretical factor of $O(1)$ and $T \equiv ke^{-\pi kR} \sim$ TeV is the mass scale for the resonances. Therefore, we want to explicitly see that these corrections do not destroy the successful relations among the electroweak gauge boson masses and couplings.

In the 4D picture, the relevant corrections of order $(c\alpha/4\pi)\ln(k/T)$ arise in the couplings of the elementary gauge fields. In a suitable normalization for the gauge bosons, they appear in the coefficients of the gauge kinetic terms for the elementary gauge bosons. In the 5D picture, this effect of running gauge couplings for the elementary gauge bosons corresponds to the radiatively-induced logarithmic energy dependence for the gauge two-point correlators whose end points are both on the Planck brane [15]. This effect is then resummed into the coefficients of the Planck-brane localized gauge kinetic terms defined at the renormalization scale of order T measured in terms of the 4D metric $\eta_{\mu\nu}$. This means that to find the corrections to the electroweak observables we simply have to calculate the masses and couplings of the γ, W and Z bosons in the presence of Planck-brane localized gauge kinetic terms which have coefficients of order $(b/8\pi^2)\ln(k/T)$, where b is the beta function coefficient. This corresponds to defining the theory at a renormalization scale of order T measured in terms of the 4D metric, integrating out the physics above this energy scale.[2]

What brane-localized kinetic terms should be included on the Planck brane at the 4D renormalization scale of order T? If the theory possesses bare Planck-brane localized gauge kinetic terms, their coefficients can take different values for $SU(2)_L$, $SU(2)_R$ and $U(1)_X$, which we denote as $1/\tilde{g}_{L,0}^2$, $1/\tilde{g}_{R,0}^2$ and $1/\tilde{g}_{X,0}^2$. In the 4D picture, these couplings correspond to tree-level gauge couplings for the elementary $SU(2)_L$, $SU(2)_R$ and $U(1)_X$ gauge fields at a scale of order k.[3] On top of this, renormalization group evolution between the energy scales of k and T induces an additional contribution to the gauge couplings. Because the elementary gauge group below the scale k is $SU(2)_L \times U(1)_Y$, this contribution will produce Planck-brane localized kinetic terms for the $SU(2)_L$ and $U(1)_Y$ gauge fields, A_μ^{La} and $A_\mu^{R3} + A_\mu^X$. We denote coefficients of these terms by $1/\tilde{g}_L^2$ and $1/\tilde{g}_Y^2$, which take the values $1/\tilde{g}_L^2 = (b_L/8\pi^2)\ln(k/T)$ and $1/\tilde{g}_Y^2 = (b_Y/8\pi^2)\ln(k/T)$ where b_L and b_Y are the beta function coefficients for $SU(2)_L$ and $U(1)_Y$. Adding all together, we obtain the Planck-brane localized kinetic terms at the scale T:

$$\begin{aligned} S \;=\; & \int d^4x \int dy\, \delta(y) \biggl[-\frac{1}{4\tilde{g}_{L,0}^2} \sum_{a=1}^{3} F_{\mu\nu}^{La} F_{\mu\nu}^{La} - \frac{1}{4\tilde{g}_{R,0}^2} \sum_{a=1}^{3} F_{\mu\nu}^{Ra} F_{\mu\nu}^{Ra} - \frac{1}{4\tilde{g}_{X,0}^2} F_{\mu\nu}^{X} F_{\mu\nu}^{X} \\ & \qquad - \frac{1}{4\tilde{g}_L^2} \sum_{a=1}^{3} F_{\mu\nu}^{La} F_{\mu\nu}^{La} - \frac{1}{16\tilde{g}_Y^2} (F_{\mu\nu}^{R3} + F_{\mu\nu}^{X})(F_{\mu\nu}^{R3} + F_{\mu\nu}^{X}) \biggr], \end{aligned} \tag{9}$$

where $F_{\mu\nu}^{R3} \equiv \partial_\mu A_\nu^{R3} - \partial_\nu A_\mu^{R3}$. The point here is that, while the bare couplings $1/\tilde{g}_{L,0}^2$, $1/\tilde{g}_{R,0}^2$ and $1/\tilde{g}_{X,0}^2$ are parameters of the theory and can take arbitrary values as long as they are positive, radiatively induced ones, $1/\tilde{g}_L^2$ and $1/\tilde{g}_Y^2$, are calculable quantities and cannot just

[2]A related calculation using the renormalization scheme adopted here can be found in [16], where various renormalization schemes and their relations are discussed in detail.

[3]In the discussions here and below, we implicitly assume the brane Higgs breaking case with large v_Σ and v_H, but essential features of our analysis are unchanged even for the boundary condition breaking case.

be made smaller. In fact, the sizes of these radiative pieces are $\sim (b/8\pi^2)\ln(k/T) = O(1)$ so we have to include the effect of these pieces. Note that, contrary to the bare couplings $1/\tilde{g}_0^2$, radiative couplings can take either negative ($1/\tilde{g}^2 < 0$) or positive ($1/\tilde{g}^2 > 0$) value, depending on whether the corresponding gauge interaction is asymptotically free or non-free. The negative $1/\tilde{g}^2$ is not problematic for the radiative piece, because this piece is vanishing at the scale k and is generated only at lower scale (in the 4D metric). At lower energies the gauge couplings squared for physical 4D gauge bosons receive contributions both from the brane piece (which could be negative) and the bulk piece (which is positive), so that the physical gauge couplings squared are always positive.

We now calculate the masses and couplings of the electroweak gauge bosons in the theory described in section 2. We neglect $SU(3)_C$ in the analysis below, since it is irrelevant for our discussion. For simplicity, we here assume that the bare brane kinetic terms are small, *i.e.* $1/\tilde{g}_{L,0}^2, 1/\tilde{g}_{R,0}^2, 1/\tilde{g}_{X,0}^2 \simeq 0$, and include only radiative pieces. In fact, this situation is naturally realized by assuming that the theory is strongly coupled at the scale k: $\tilde{g}_{L,0} \sim \tilde{g}_{R,0} \sim \tilde{g}_{X,0} \sim 4\pi$. The effects of introducing bare couplings are discussed at the end of this section.

We first solve mass eigenvalues of the W and Z gauge-boson towers. The gauge kinetic terms are given by Eqs. (2, 9), and the boundary conditions for the gauge fields are given by Eqs. (3, 4) appropriately modified by the presence of the brane kinetic terms. In particular, we derive the masses of the lowest towers of W and Z, which we identify as the standard-model electroweak gauge bosons. Considering the sizes of various couplings, $1/g_L^2 \sim 1/g_R^2 \sim 1/g_X^2 \sim 1/\pi R$ and $1/\tilde{g}_L^2 \sim 1/\tilde{g}_Y^2 \sim 1$,[4] we obtain the following expressions for the W and Z boson masses:

$$m_W^2 = \frac{1}{\frac{\pi R}{g_L^2} + \frac{1}{\tilde{g}_L^2}} \frac{2\,T^2}{(g_L^2 + g_R^2)k}, \tag{10}$$

$$m_Z^2 = \frac{\frac{\pi R}{g_L^2} + \frac{\pi R}{g_R^2} + \frac{\pi R}{g_X^2} + \frac{1}{\tilde{g}_L^2} + \frac{1}{\tilde{g}_Y^2}}{\left(\frac{\pi R}{g_L^2} + \frac{1}{\tilde{g}_L^2}\right)\left(\frac{\pi R}{g_R^2} + \frac{\pi R}{g_X^2} + \frac{1}{\tilde{g}_Y^2}\right)} \frac{2\,T^2}{(g_L^2 + g_R^2)k}, \tag{11}$$

at the leading order in T/k and in $1/\pi kR$. We have checked numerically that these expressions approximate the true values at the level of a percent, which is sufficient for our purpose. The photon mass is zero, $m_\gamma = 0$, as it should be.

The couplings of the electroweak gauge bosons to the quarks and leptons can be read off from Eqs. (7, 8). By normalizing wavefunctions such that the kinetic terms for the 4D gauge fields are canonically normalized, *i.e.* $\mathcal{L}_{\rm 4D} = -(1/4)F_{\mu\nu}F_{\mu\nu}$, these equations give fermion couplings

[4] Because the 4D gauge couplings, $g_{\rm 4D}$, are $O(1)$ and related to the bulk gauge couplings, g_B, by $1/g_{\rm 4D}^2 \sim \pi R/g_B^2$, g_B's should be regarded as quantities of $O(\sqrt{\pi R})$.

to γ, W and Z bosons. We define the fermion gauge coupling parameters g_W, e, g_Z and η_Z as

$$\begin{aligned} \mathcal{D}_\mu \psi_L &= \partial_\mu \psi_L + \frac{i}{2} g_W \begin{pmatrix} 0 & W_\mu^1 - iW_\mu^2 \\ W_\mu^1 + iW_\mu^2 & 0 \end{pmatrix} \psi_L \\ &\quad + ie \begin{pmatrix} Q_{\mathrm{EM},L\uparrow} & 0 \\ 0 & Q_{\mathrm{EM},L\downarrow} \end{pmatrix} A_\mu^\gamma \psi_L + i g_Z \begin{pmatrix} -Y + \frac{1}{2}\eta_Z & 0 \\ 0 & -Y - \frac{1}{2}\eta_Z \end{pmatrix} Z_\mu \psi_L \end{aligned} \qquad (12)$$

for left-handed fermions $\psi_L = q, l$, and

$$\mathcal{D}_\mu \psi_R = \partial_\mu \psi_R + ie Q_{\mathrm{EM},R} A_\mu^\gamma \psi_R + i g_Z (-Y) Z_\mu \psi_R \qquad (13)$$

for right-handed fermions $\psi_R = \bar{u}, \bar{d}, \bar{e}$. Here, W_μ, A_μ^γ and Z_μ are the canonically normalized 4D electroweak gauge bosons, and Q_{EM}'s and Y are the electric charges and hypercharges. In the present theory, we find that correct values for the electric charges and hypercharges, $Q_{\mathrm{EM},L\uparrow}$, $Q_{\mathrm{EM},L\downarrow}$, $Q_{\mathrm{EM},R}$ and Y, are reproduced for q, $\bar{u}$, $\bar{d}$, l and $\bar{e}$ (and $\bar{n}$ as well) with the quarks and leptons in the representations of Eq. (6). The couplings g_W, e, g_Z and η_Z are given by

$$\frac{1}{g_W^2} \simeq \frac{\pi R}{g_L^2} + \frac{1}{\tilde{g}_L^2}, \qquad (14)$$

$$\frac{1}{e^2} \simeq \frac{\pi R}{g_L^2} + \frac{\pi R}{g_R^2} + \frac{\pi R}{g_X^2} + \frac{1}{\tilde{g}_L^2} + \frac{1}{\tilde{g}_Y^2}, \qquad (15)$$

$$\frac{1}{g_Z^2} \simeq \left(\frac{\pi R}{g_R^2} + \frac{\pi R}{g_X^2} + \frac{1}{\tilde{g}_Y^2} \right) + \left(\frac{\pi R}{g_L^2} + \frac{1}{\tilde{g}_L^2} \right)^{-1} \left(\frac{\pi R}{g_R^2} + \frac{\pi R}{g_X^2} + \frac{1}{\tilde{g}_Y^2} \right)^2, \qquad (16)$$

$$\frac{1}{\eta_Z} \simeq \left(\frac{\pi R}{g_L^2} + \frac{1}{\tilde{g}_L^2} \right) \left(\frac{\pi R}{g_R^2} + \frac{\pi R}{g_X^2} + \frac{1}{\tilde{g}_Y^2} \right)^{-1}, \qquad (17)$$

at the leading order in T/k and in $1/\pi k R$, and there is no other couplings than those of the form given in Eqs. (12, 13).

Now, we compare our results, Eqs. (10, 11, 14, 15, 16, 17), with the corresponding tree-level expressions in the standard model. In the standard model, the masses and couplings of the electroweak gauge bosons are given by

$$m_W^2 = \frac{g^2}{2} v^2, \qquad m_Z^2 = \frac{g^2 + g'^2}{2} v^2, \qquad (18)$$

and

$$\frac{1}{g_W^2} = \frac{1}{g^2}, \qquad \frac{1}{e^2} = \frac{1}{g^2} + \frac{1}{g'^2}, \qquad \frac{1}{g_Z^2} = \frac{1}{g'^2} + \frac{g^2}{g'^4}, \qquad \frac{1}{\eta_Z} = \frac{g'^2}{g^2}, \qquad (19)$$

respectively, where g and g' represent the $SU(2)_L$ and $U(1)_Y$ gauge couplings and v is the VEV of the Higgs field. From these expressions we find that the present theory completely reproduce the tree-level structure of the standard model at the level considered here. The correspondence between the two theories is given by

$$\frac{1}{g^2} = \frac{\pi R}{g_L^2} + \frac{1}{\tilde{g}_L^2}, \qquad \frac{1}{g'^2} = \frac{\pi R}{g_R^2} + \frac{\pi R}{g_X^2} + \frac{1}{\tilde{g}_Y^2}, \qquad v^2 = \frac{4\,T^2}{(g_L^2 + g_R^2)k}. \tag{20}$$

This generalizes the result of [6] to the case with the arbitrary brane kinetic terms and the bulk gauge couplings, especially to $g_L \neq g_R$.

Note that the brane pieces $1/\tilde{g}_{L,Y}^2$ cannot just be neglected because they are calculable and have size of order $(b/8\pi^2)\ln(k/T) \sim 1$. In fact, assuming that the theory has the minimal matter content, *i.e.* three generations of quarks and leptons with the quantum numbers given in Eq. (6), and unwanted fields are made heavy by Planck-brane masses of order k, the running of the brane gauge couplings $\tilde{g}_L$ and $\tilde{g}_Y$ is given by that of the standard model without a Higgs boson. The brane kinetic terms at the scale T are then given by

$$\frac{1}{\tilde{g}_L^2} = \frac{b_L}{8\pi^2}\ln\Bigl(\frac{k}{T}\Bigr) \simeq -1.26, \qquad \frac{1}{\tilde{g}_Y^2} = \frac{b_Y}{8\pi^2}\ln\Bigl(\frac{k}{T}\Bigr) \simeq 2.53, \tag{21}$$

where $(b_L, b_Y) = (-10/3, 20/3)$ and we have taken $\ln(k/T) \simeq 30$. Therefore, any computation in the present theory must include the effect of these terms.

The effect of bare Planck-brane gauge couplings can be included by replacing $\tilde{g}_L^2$ and $\tilde{g}_Y^2$ in Eqs. (10, 11, 14, 15, 16, 17) as

$$\frac{1}{\tilde{g}_L^2} \;\rightarrow\; \frac{1}{\tilde{g}_L'^2} = \frac{1}{\tilde{g}_{L,0}^2} + \frac{1}{\tilde{g}_L^2}, \qquad\qquad \frac{1}{\tilde{g}_Y^2} \;\rightarrow\; \frac{1}{\tilde{g}_Y'^2} = \frac{1}{\tilde{g}_{R,0}^2} + \frac{1}{\tilde{g}_{X,0}^2} + \frac{1}{\tilde{g}_Y^2}. \tag{22}$$

Therefore, it does not affect the masses and couplings of the electroweak gauge bosons either, at the leading order in T/k and $1/\pi kR$. Incidentally, introduction of bare brane couplings does not allow to have $1/\tilde{g}_L'^2 \simeq 1/\tilde{g}_Y'^2 \simeq 0$ in the simplest setup considered here, because Eq. (21) and positive bare couplings, $1/\tilde{g}_{R,0}^2, 1/\tilde{g}_{X,0}^2 > 0$, imply $1/\tilde{g}_Y'^2 \gtrsim 2.5$. Note that, contrary to the radiative ones, bare couplings cannot be negative, since they represent physical gauge couplings for the elementary gauge bosons at high energies so that negative bare couplings lead to the appearance of ghosts in the physical states.

The correspondence obtained in Eq. (20) has some important physical consequences. Suppose we numerically calculate the deviations from the standard model relations by exactly solving the masses and couplings of W, Z and γ. We then find that they are suppressed by $1/\pi kR$ and thus at a level of a percent as claimed in [6]. Moreover, we find that with the observed 4D couplings fixed by Eq. (20), these $1/\pi kR$ deviations become smaller when g_R/g_L becomes larger. This is because in this parameter region the value of T becomes larger and thus the masses of the resonances become higher (see the discussion around Eq. (27) below). This is an important result because it allows us to control the scale of new physics in theories with dynamical electroweak symmetry breaking, by changing the parameters of the theory.

While we do not attempt to perform a full analysis of higher order effects, such as $1/\pi kR$ suppressed corrections and matter loop effects, we think it is significant that there is a theory reproducing the standard model predictions at a level of a percent or less. Because the experimental data require deviations from the standard model relations to be smaller than a factor of a few times 10^{-3} for the electroweak gauge boson masses and couplings, it will still be needed to suppress the deviations by a factor of a few to an order of magnitude. These may be attained by pushing up the scale of T by changing g_R/g_L and/or by cancellations among various contributions such as those from Planck-brane gauge couplings, calculable radiative corrections and TeV-brane localized gauge couplings. Although the situation is not completely clear, in the next section we assume that these are somehow attained and present a (potentially) realistic theory of electroweak symmetry breaking without a Higgs boson. We find that the theory can reproduce many successful features of the standard model, including fermion mass generation, while satisfying various constraints from experiments. We also discuss phenomenological implications of the theory such as small non-universality of fermion couplings to the W and Z bosons and the spectrum of the gauge-boson KK resonances.

4 Potentially Realistic Theory and Its Implications

In this section we present a theory of electroweak symmetry breaking without a Higgs boson, along the line of section 2. In the 4D dual picture, the theory below the scale k appears as the standard model without a Higgs field. It also contains an extra gauge group G which becomes strong at the electroweak scale, breaking the electroweak symmetry. We present a complete model including the quark and lepton sector which reproduces the observed structure for the quark and lepton masses and couplings. Some elements for constructing our theory can be found in the literatures [6, 10, 7, 13, 14]; in particular, the basic gauge structure of the theory has been given in [6]. We also discuss several phenomenological implications of the theory.

As already discussed in section 2, we consider an $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_X$ gauge theory in 5D warped space, with the metric given in Eq. (1). We here consider breaking these groups by Higgs fields localized at the Planck and TeV branes, whose representations and VEVs have been given in the paragraph containing Eq. (5). In the limit that v_Σ and v_H are large, this breaking becomes equivalent to imposing boundary conditions in Eqs. (3, 4); in particular, the physical Higgs bosons decouple from low-energy physics and the theory does not contain any narrow scalar resonance. At low energies unbroken gauge symmetry of the theory is $SU(3)_C \times U(1)_{\rm EM}$, and the weak interaction is mediated by the massive W and Z bosons.

Three generations of quarks and leptons are introduced in the bulk as discussed in section 2 [in the paragraph containing Eq. (6)]. Specifically, a single generation of quarks and leptons are contained in $q(\mathbf{3},\mathbf{2},\mathbf{1},1/6)$, $\psi_{\bar{u}}(\mathbf{3}^*,\mathbf{1},\mathbf{2},-1/6)$, $\psi_{\bar{d}}(\mathbf{3}^*,\mathbf{1},\mathbf{2},-1/6)$, $l(\mathbf{1},\mathbf{2},\mathbf{1},-1/2)$ and $\psi_{\bar{e}}(\mathbf{1},\mathbf{1},\mathbf{2},1/2)$ with the unwanted components made heavy by the orbifold compactification and mixing with fields localized on the Planck brane. This reproduces the standard model matter content at low energies. The fermion masses arise from the couplings introduced on the TeV brane

$$S = \int d^4x \int dy\; \delta(y-\pi R)\sqrt{-g_{\rm ind}}\Big[y_u q \psi_{\bar{u}} H + y_d q \psi_{\bar{d}} H + y_e l \psi_{\bar{e}} H + {\rm h.c.} \Big], \tag{23}$$

where we have suppressed the generation index, $i = 1,2,3$. As the up-type quark, down-type quark, and charged lepton masses arise from three independent couplings, y_u, y_d and y_e, it is clear that there are no unwanted relations among them coming from $SU(2)_R$. Small neutrino masses can be naturally obtained through the seesaw mechanism, by introducing right-handed neutrino fields $\psi_{\bar{n}}(\mathbf{1},\mathbf{1},\mathbf{2},1/2)$ in the bulk together with the TeV-brane couplings $\delta(y-\pi R) y_n l \psi_{\bar{n}} H$ and the Planck-brane Majorana masses $\delta(y)(\psi_{\bar{n}}\Sigma^\dagger)^2$.[5] We choose the bulk masses for the first-two generation fermions to be larger than $k/2$, *i.e.* $c > 1/2$ [for the definition of c's, see section 2]. For $c > 1/2$, the zero modes for the fermions have the wavefunctions localized to the Planck brane, so that the 4D masses for these quarks and leptons are naturally suppressed. This wavefunction localization also suppresses effects from unwanted flavor violating operators on the TeV brane, such as $q_i^\dagger q_j q_k^\dagger q_l$.[6] The bulk masses for the third generation q and $\psi_{\bar{d}}$ can be chosen to be around $c = 1/2$. To obtain the top quark mass we consider the right-handed top quark to have $c < 1/2$. Note that in our theory the value of T is not strictly related to the W and Z boson masses [see Eqs. (25, 26) below], so that we can have larger values of T by choosing parameters of the model. Successful top phenomenology will require the value of T close to its maximum, $T \simeq 900$ GeV, but we leave a detailed study of this issue for future work. The value of T also becomes important when we discuss KK towers of the gauge bosons.

We now consider fermion couplings to the electroweak gauge bosons. As we have seen in the previous section, these couplings are those of the standard model for Planck-brane localized matter, at the leading order in $1/\pi kR$. We here assume that this is the case at all orders for

[5]The operators $\delta(y-\pi R) l^2 H^2$ and $\delta(y-\pi R) l^2 H^{\dagger 2}$, which could potentially give large masses for the observed neutrinos, are forbidden by gauge invariance.

[6]The amount of localization consistent with the fermion masses is not enough to suppress proton decay caused by operators on the TeV brane, but this can easily be forbidden by imposing the baryon number symmetry on the theory.

the Planck-brane matter and consider the effects of delocalizing matter in the fifth dimension. Since the W and Z boson wavefunctions have non-trivial profiles in the extra dimension, this introduces flavor non-universality in the gauge boson couplings to matter and could potentially be dangerous. We first consider the W boson couplings to the left-handed fields q and l. These couplings are obtained by convolving the wavefunction for q (or l) with that of the W boson and thus depend on the bulk mass for the fermion field. However, we find that the resulting non-universality is highly suppressed for c larger than $1/2$. It is less then 10^{-4} for $c \gtrsim 0.6$ and less than a percent even for $c = 0.5$. This effect, therefore, does not give a contradiction to experimental data. The delocalization also introduces the couplings of the right-handed fields, $\bar{u}$, $\bar{d}$ and $\bar{e}$, to W, and these couplings have a similar size to the non-universality discussed above [10^{-4} for $c \gtrsim 0.6$ and at a level of a percent for $c = 0.5$]. They, however, do not connect between the known particles, say $\bar{u}$ and $\bar{d}$, but connect the standard-model particles with some heavy fields, say $\bar{u}$ with the other component of $\psi_{\bar{u}}$. Therefore, this effect is also negligible.

How about the non-universality of the Z couplings? This is potentially dangerous because it induces FCNCs. However, we again find that the deviation from the universality is small: $\lesssim 10^{-4}$ for $c \gtrsim 0.6$, and even for $c \simeq 1/2$ it is less than a percent for left-handed fermions and at most a few percents for right-handed fermions. Although this level of non-universality could still lead to FCNCs at a level of the experimental bounds, here we stick to rough estimates and consider that the theory can evade the bounds by having c's larger than about 0.6 for the first-two generation fermions.[7] The effect could be larger for the third generation fields because of the smaller values for c. This may have observable consequences, but we do not perform a detailed analysis here for these processes. FCNCs generated through the KK resonances are expected to have a similar size. Incidentally, the fermion couplings to the photon are completely universal due to the unbroken $U(1)_{\rm EM}$ gauge invariance (because the wavefunction for the photon is flat), so no experimental constraint arises from these couplings.[8]

We finally discuss KK towers for the gauge bosons. The masses for these towers are given, regardless of the values for the bulk and brane gauge couplings, by

$$m_n \simeq \frac{\pi}{2}\Big(n + \frac{1}{2}\Big)T, \tag{24}$$

where $n = 1, 2, 3, 4 \cdots$ for the W and Z towers and $n = 1, 3, 5, \cdots$ for the γ tower (the gluon tower has the same spectrum as the γ tower). What is the value of T in the present theory?

[7] We can always choose c's for the first-two generation fermions with the same gauge quantum numbers to be equal (or very close) to evade the bounds from FCNCs. The fermion mass hierarchy is then generated by the structure of the TeV-brane couplings, y_u, y_d and y_e.

[8] The ratios between the triple gauge-boson vertex γWW and the fermion couplings to γ are exactly those of the standard model for the same reason, while the vertex ZWW receives corrections of $O(1\%)$.

From Eqs. (10, 11) we find that it is related to the masses of the W and Z bosons. This relation, however, does not uniquely determine the value of T with the given W and Z masses. Neglecting small corrections, the relation can be written as

$$T^2 \simeq \frac{\pi k R}{2 g_W^2} \frac{g_L^2}{\pi R} \left(1 + \frac{g_R^2}{g_L^2} \right) m_W^2, \tag{25}$$

so we find that T depends on the ratio of the bulk gauge couplings g_R^2/g_L^2. The values of g_L and g_R, however, are constrained to reproduce the observed W and Z boson couplings. Thus, with the given brane kinetic terms, we can determine the allowed region for T and consequently the masses for the KK towers.

As an example, let us consider the case with the minimal matter content. In this case, the brane kinetic terms are given by $1/\tilde{g}_L^2 \simeq -1.26$ and $1/\tilde{g}_Y^2 \simeq 2.53$ (see Eq. (21)),[9] so that the bulk gauge couplings satisfy the relations $\pi R/g_L^2 = 1/g_W^2 - 1/\tilde{g}_L^2 \simeq 3.61$ and $\pi R/g_R^2 + \pi R/g_X^2 = 1/e^2 - 1/g_W^2 - 1/\tilde{g}_Y^2 \simeq 5.32$. Here, we have used the experimentally measured values of $g_W \simeq 0.652$ and $e \simeq 0.313$. While $\pi R/g_L^2$ is uniquely determined, $\pi R/g_R^2$ is not, although the latter relation implies $\pi R/g_R^2 \lesssim 5.32$. The lower bound on the value of $\pi R/g_R^2$ then comes from the requirement that the 5D theory is not strongly coupled up to the cutoff scale, M_*, of the theory. This gives $\pi R/g_R^2 \gtrsim c\pi M_* R/24\pi^3 \approx 0.1 M_*/k$. Requiring $M_*/k \gtrsim 3$ so that the AdS solution is trustable, we find $0.7 \lesssim g_R^2/g_L^2 \lesssim 12$. Using this in Eq. (25), together with $\pi k R \simeq 30$ and $m_W \simeq 80$ GeV, we obtain

$$330 \text{ GeV} \lesssim T \lesssim 900 \text{ GeV}, \tag{26}$$

which is translated into the values for the masses of the first KK towers of γ, W and Z (and gluon):

$$800 \text{ GeV} \lesssim m_1 \lesssim 2.1 \text{ TeV}. \tag{27}$$

Here, smaller (larger) values for m_1 correspond to smaller (larger) values for g_R^2/g_L^2. Although the couplings of these excitations to the matter fields are suppressed by small wavefunction overlaps for $c > 1/2$, the above mass region is still within the reach of the LHC [17]. The flexibility for the masses of the KK towers also helps to control the size of corrections to the electroweak observables. In fact, for $g_R \gg g_L$, the theory does not have any new particle than the standard model gauge bosons and quarks and leptons up to an energy scale of 2 TeV, except possible states associated with the top quark t'. In particular, no physical Higgs boson appears

[9] Having $c < 1/2$ for the right-handed top quark may give somewhat smaller values for $1/\tilde{g}_Y^2$ because the value of b_Y becomes smaller depending on the value of c for $c < 1/2$. The resulting correction, however, is small.

below, and even above, this energy scale.[10] It will be interesting to see in detail whether theories such as the one described here can really have acceptable parameter regions which comply with all the experimental data including the precision electroweak data.

5 Conclusions

We have considered a theory of electroweak symmetry breaking without a Higgs boson. The theory is formulated in 5D warped space, but through AdS/CFT duality, it is related to the 4D theory in which the electroweak symmetry is dynamically broken by non-perturbative effects of some gauge interaction G. We have studied several issues in this framework, including masses and couplings of the electroweak gauge bosons and the structure for the quark and lepton mass generation.

We have studied radiative corrections to the masses and couplings of the electroweak gauge bosons in 5D which are logarithmically enhanced by the mass ratio between the Planck and TeV scales. We explicitly showed that these corrections do not destroy the successful relations among the masses and couplings of W, Z and γ at the leading order. We have also extended the previous result in [6] to the case of most general brane and bulk gauge couplings. This allows us to incorporate quarks and leptons in a relatively simple manner and also leads to an extra free parameter in the relation between the scale of new physics and the W and Z boson masses. The latter point seems particularly interesting as it allows us to push up the masses of the resonances and suppress the deviations from the standard model relations arising from loops of these states. Although it is still not completely clear whether the theory can naturally evade all the precision constraints, it is certainly encouraging if we could have a theory without a Higgs boson which successfully reproduces the standard model structure at a level of a percent or less. We have presented a concrete such model which can accommodate realistic quark and lepton masses and mixings as well as their couplings to the electroweak gauge bosons. An interesting possibility is that the observed quark and lepton mass structure is obtained through

[10] In spite of the absence of a physical Higgs boson, unitarity of the theory is maintained up to a high energy of order the 4D Planck scale. This is clear in the 5D picture because the theory is a well-defined effective field theory below the cutoff scale M_*, which appears close to the 4D Planck scale for an observer on the Planck brane. From the low-energy bottom-up point of view, bad high energy behaviors of the longitudinal WW scattering are unitarized by the appearance of the KK states and the effective vertices arising by integrating out these states [10, 6, 18]. This will make the theory to be unitary up to the scale around 10 TeV, above which the simple KK picture is not necessarily appropriate. However, the theory above this scale is simply an unbroken $SU(3)_C \times SU(2)_L \times U(1)_Y \times G$ gauge theory so it is unitary up to very high energies. In fact, this unitarity consideration suggests that the first KK excitation of the electroweak gauge bosons appears below ~ 2 TeV, where the E^2 amplitude of the longitudinal WW scattering would blow up. It is interesting that the maximum value of the lowest KK mass obtained in Eq. (27) saturates this bound.

wavefunction profiles of these fields in the extra dimension. This leads to some amount of flavor non-universality in the matter couplings to the gauge bosons. We have estimated this effect and found that it does not lead to an immediate contradiction with the experimental bounds. We have also considered the KK excited states for the gauge bosons and found that they are within the reach of the LHC.

We have not performed a full analysis for the corrections to the electroweak observables at a level of order a percent. These corrections come from tree-level brane kinetic terms at the Planck and TeV branes and calculable radiative corrections from gauge and matter loops. Because the corrections of this order are still too large to satisfy constraints from the precision electroweak data, they have to be suppressed by a factor of a few to an order of magnitude. These may be attained by pushing up the mass scale of the resonances and/or by cancellations among various contributions. It may also be possible for these corrections to mimic a light Higgs boson suggested by the precision electroweak data, by appropriately choosing parameters of the theory. It will be interesting to study these corrections in more detail.

Acknowledgments

I would like to thank David Smith for discussions. I also thank Kaustubh Agashe and Alex Pomarol for useful discussions which have helped me to find and correct an error in the earlier version of this paper. Fermilab is operated by Universities Research Association, Inc. for the U.S. Department of Energy under contract DE-AC02-76CH03000.

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