The Gaugephobic Higgs

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Abstract

We present a class of models that contains Randall-Sundrum and Higgsless models as limiting cases. Over a wide range of the parameter space WW scattering is mainly unitarized by Kaluza-Klein partners of the W and Z, and the Higgs particle has suppressed couplings to the gauge bosons. Such a gaugephobic Higgs can be significantly lighter than the 114 GeV LEP bound for a standard Higgs, or heavier than the theoretical upper bound. These models predict a suppressed single top production rate and unconventional Higgs phenomenology at the LHC: the Higgs production rates will be suppressed and the Higgs branching fractions modified. However, the more difficult the Higgs search at the LHC is, the easier the search for other light resonances (like Z', W', t', exotic fermions) will be.

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1 Introduction

Extra-dimensions provide a natural explanation of the hierarchy between the Planck and the electroweak scale. In Randall-Sundrum (RS) scenarios [1] the hierarchy is explained through an exponential warping of mass scales along an extra-dimension. Early versions of these scenarios had all the Standard Model (SM) fields localized on the TeV brane [1]. Later the gauge fields migrated to the bulk [2] and were soon followed by the quarks and leptons [3,4]. This migration made it easier to understand how the scenario could be compatible with unification and precision electroweak measurements even with TeV scale Kaluza-Klein (KK) modes for these fields. However the Higgs remained mainly[1] on the brane, since this was the main motivation of the RS scenario: to explain how the Higgs could naturally have a TeV scale mass. However from the perspective of the anti-de Sitter/conformal field theory (AdS/CFT) correspondence, the localization of the Higgs on the TeV brane corresponds to the assumption that the operator that breaks the electroweak symmetry has an infinite scaling dimension. Furthermore from the AdS/CFT perspective the technical naturalness of the Higgs mass only requires that the operator have a scaling dimension greater than two so that the operator corresponding to the Higgs squared mass term has a dimension greater than four and is thus irrelevant. Therefore it is interesting to explore what happens in RS scenarios when the assumption of strict localization of the Higgs field is relaxed.

Another generalization of RS scenarios was the suggestion [5] that with an extra dimension, some or all of the W and Z masses could come from the momentum along the extra dimension, that is from the fact that their wave functions are not exactly flat. By *increasing* the Higgs VEV on the brane the lightest W and Z modes are forced to move further from the brane, and their wave functions become less zero-mode like. In the limit that the brane Higgs VEV goes to infinity the W and Z have no couplings to the Higgs and all of their masses come from the derivative of their extra dimensional profiles. This is the Higgsless limit [5]. One can then interpolate between the Higgsless model, the conventional RS model and the SM by varying the Higgs VEV and the Higgs localization parameter, and as a result also varying the size of the coupling of the Higgs to the gauge bosons. Since the integral of the Higgs VEV contributes to the gauge boson masses we cannot take the infinite VEV limit unless the Higgs is strictly localized on the brane. Thus these bulk Higgs theories generically have a physical Higgs with suppressed gauge couplings. In conventional RS scenarios the Higgs VEV is tuned to be much smaller than the typical scale of the model (i.e. the scale of the KK modes). Again from the AdS/CFT perspective this is not the natural expectation. Normally we would expect at most a factor of 4π between the analogs of the pion decay constant and the rho mass.[2] Of course, at the LHC we would like to be prepared for the unexpected, so when considering RS scenarios we should keep in mind that there is a two dimensional parameter space (Higgs localization and VEV) to consider, and that the most "natural" locations in the parameter space are not those that have been studied so far. Due to the repulsion of the wave functions of the bulk fields away from the Higgs VEV we find

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[1]See however refs. [6,7].

[2]A natural way of keeping the Higgs VEV below the KK scale is via the pseudo-Goldstone mechanism. A concrete realization of this idea in the RS context has been proposed in ref. [8].

that all Higgs couplings are generically suppressed relative to the SM, and this has a variety of consequences for experimental searches.

In this paper we will attempt a first survey of this gaugephobic Higgs parameter space. In Section 2 we describe the class of models (with many of the details relegated to the three appendices) and how they are consistent with current experiments, while in Section 3 we describe the phenomenology for upcoming experiments. To keep the discussion simple we focus on two benchmark points: the first where the Higgs couplings are about half of the SM values, and a second where the couplings are about a tenth of the SM values. The Higgs is not only harder to find, but it can also be outside the range of masses allowed in the SM: it can be lighter than the experimental bound from LEP or heavier than the theoretical upper bound coming from unitarity. Fortunately there are a variety of other particles that can be searched for but typically it will take much longer for the LHC to sort out the true situation than in other more simplistic scenarios.

2 The model

In this section we examine the effect of a bulk Higgs in an RS set-up using the conformally flat metric

$$ds^2 = \left(\frac{R}{z}\right)^2 \left(\eta_{\mu\nu} dx^\mu dx^\nu - dz^2\right) , \tag{2.1}$$

for $R < z < R'$. As usual the UV brane is located at $z = R$ and the TeV brane is at $z = R'$. Here we will simply adapt the Higgsless model of ref. [5, 9] by replacing the Dirichlet IR boundary conditions (BC's) (corresponding to a very large Higgs VEV localized on the IR brane) with a bulk Higgs VEV. Let us first summarize the main features of the model: the gauge group is $SU(2)_L \times SU(2)_R \times U(1)_X$, where for the first two generations the X charge is equivalent to $B - L$. The gauge bosons for each group are denoted by $A_{L,R}^a$ and B, with 5D gauge couplings g_5 and $\tilde{g}_5$ (for simplicity we will take the two $SU(2)$'s to have the same gauge coupling). The Higgs field is a bidoublet of $SU(2)_L \times SU(2)_R$:

$$\mathcal{H} = \begin{pmatrix} \phi_0^* & \phi_+ \\ -\phi_+^* & \phi_0 \end{pmatrix} \qquad \mathcal{H} \to U_L \mathcal{H} U_R^\dagger ; \tag{2.2}$$

and its $U(1)_X$ charge is zero. The Lagrangian for the Higgs and the $SU(2)_L \times SU(2)_R \times U(1)_X$ gauge sector is:

$$
\begin{aligned}
\mathcal{L} = {} & \int_R^{R'} dz \frac{R}{z} \left\{ -\frac{1}{4g_5^2} F_L^{aMN} F_{LMN}^a - \frac{1}{4g_5^2} F_R^{aMN} F_{RMN}^a - \frac{1}{4\tilde{g}_5^2} B^{MN} B_{MN} \right\} \\
& + \int_R^{R'} dz \left(\frac{R}{z}\right)^3 \left[\mathrm{Tr}\, |\mathcal{D}_M \mathcal{H}|^2 - \frac{\mu^2}{z^2} \mathrm{Tr}\, |\mathcal{H}|^2 \right] - V_{\mathrm{UV}}(\mathcal{H})\delta(z - R) - V_{\mathrm{TeV}}(\mathcal{H})\delta(z - R') ,
\end{aligned}
\tag{2.3}
$$

where μ is a bulk mass for the Higgs (in units of the inverse curvature radius R^{-1}). The potentials V_{UV} and V_{TeV} on the branes determine the boundary conditions for the Higgs

and induce electroweak symmetry breaking (EWSB). In particular the potential on the TeV brane which breaks the electroweak symmetry has the usual form:

$$V_{\text{TeV}} = \left(\frac{R}{R'}\right)^4 \frac{\lambda R^2}{2} \left(\text{Tr}\,|\mathcal{H}|^2 - \frac{v_{\text{TeV}}^2}{2}\right)^2 . \tag{2.4}$$

The bulk equations of motion allow two solutions for the profile of the bulk Higgs VEV (see appendix A). We choose a UV boundary condition that ensures that the allowed solution is localized close to the TeV brane (i.e. electroweak symmetry is broken at the TeV scale). The solution for the bulk profile can be written as:

$$v(z) = \sqrt{\frac{2(1+\beta)\log R'/R}{(1-(R/R')^{2+2\beta})}} \frac{gV}{g_5} \frac{R'}{R} \left(\frac{z}{R'}\right)^{2+\beta} , \tag{2.5}$$

where g is the SM $SU(2)$ gauge coupling, and

$$\beta = \sqrt{4 + \mu^2} \tag{2.6}$$

is the parameter determining how close to the TeV brane the Higgs VEV is localized. In the presence of this VEV the gauge boson masses have two sources: the curvature of their wave functions (also induced by the VEV) and the direct overlap with the VEV. Instead of using v_{TeV} of Eq. (2.4), we choose to normalize the VEV through the input parameter V (appearing in Eq. (2.5), which carries the usual dimension of mass). The factors of β, R and R' in front are chosen in such a way that in the limit $V = v_{\text{SM}} \sim 246$ GeV the direct contribution of the VEV to the gauge bosons mass saturates the SM value, and the volume of the extra dimension shrinks to zero. In other words one recovers the SM in 4 dimensions, independently of the localization of the VEV.

The theory also contains a bulk physical Higgs boson. The mass of the (possibly) light mode is a free parameter and determined by the quartic coupling λ of the TeV potential V_{TeV}. Similarly to the SM the mass of the physical Higgs is given by

$$m_h^2 \propto \lambda V^2 . \tag{2.7}$$

It is clear that increasing V the Higgs can be decoupled safely from the theory without entering a strongly coupled regime in the Higgs sector, as happens in the SM.

The model has 6 free parameters in the gauge-Higgs sector: R, R', the gauge couplings g_5 and $\tilde{g}_5$, the VEV parameters β and V. Three of the parameters (R', g_5 and $\tilde{g}_5$) can be determined by fixing the values of the W and Z mass, and the coupling of the photon to the fermions. If we also fix $1/R = 10^8$ GeV [1], we are left with the two parameters (V, β). In Fig. 1 we plot contours of the IR scale $1/R'$ as a function of V and β. It is interesting to notice various limits. First, when the VEV V is equal to the SM value $v_{\text{SM}} \sim 246$ GeV, the IR scale becomes large. In fact, in this case, the contribution of the Higgs VEV to the

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[1]Different values of R will not qualitatively affect our results. The main effect of R is to enter in the relation between M_W and R': for instance, a smaller R will require a larger scale on the TeV brane $1/R'$ [10].
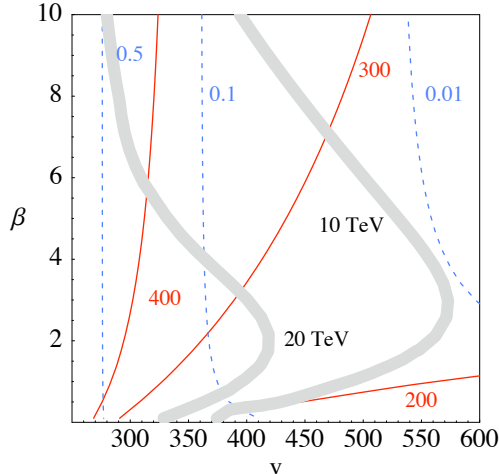
Figure 1: Lines with fixed values (400 Gev, 300 GeV, and 200 GeV) of the inverse size of the extra-dimension R' in the (V, β) plane (continuous lines), lines of fixed contribution of the Higgs to WW scattering compared to the SM (dashed lines), and lines of fixed cut-off of the theory due to the top sector becoming non-perturbative (thick gray lines).

W mass saturates the physical value of it, and thus the contribution of the extra dimension has to vanish. As a consequence, in this limit (with fixed W mass), the volume of the extra dimensions has to go to zero and the KK modes decouple: $R' \to R$. In other words, when $V \to v_{SM}$ we recover the 4D SM.

Another interesting limit is $\beta \to \infty$. This is a conventional RS1 scenario [11]. If one also takes the limit $V \gg 1/R'$ (i.e. top-right corner of the plot), one recovers the Higgsless limit where the Higgs decouples from the theory and can be made arbitrarily heavy: the unitarization of the longitudinal W scattering is then guaranteed by the contribution of the gauge boson resonances [12, 13]. In the bulk of the parameter space, however, we find an intermediate situation where both the Higgs and the gauge bosons KK modes contribute. In the SM, the Higgs is forced to be lighter than about 1 TeV in order to play a role in the unitarization (unitarity bound). As a consequence, the SM Higgs cannot escape detection at the LHC. In the Higgsless limit, the Higgs cannot be detected in any experiment. In the intermediate regime, the role of the Higgs in the unitarization can be less important than in the SM, and the usual Higgs mass bound can be evaded. In cases where the Higgs is out of reach of the LHC, the world will appear to be Higgsless from the experimental point of view for several decades. It is important then to understand the unitarity bound in this model. The Higgs enters in the longitudinal W scattering amplitude term which grows like the energy squared. The coefficient of this term is proportional to

$$\mathcal{A}^{(2)} \sim g_{WWWW}^2 - \frac{3}{4} \sum_k \frac{M_{Z^k}^2}{M_W^2} g_{WWZ^k}^2 - \frac{1}{4} \sum_k g_{WWH^k}^2 \,, \tag{2.8}$$

where the sums cover all the KK modes. In the SM, only the first Z and Higgs mode are

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present and both sums consist of only one term. The presence of the gauge boson KK modes increases the second term, allowing the third one (the Higgs contribution) to be smaller. We can thus define the following parameter to quantify how "Higgsless" the model is:

$$\xi \equiv \frac{\sum_k g_{WWH^k}^2}{g_{WWH}^2(SM)}. \tag{2.9}$$

In the SM limit $\xi \to 1$, while in the Higgsless limit $\xi \to 0$. Since 5D gauge invariance guarantees the vanishing of (2.8) this implies that the contribution of the lightest Higgs will not be important until scales of order $\Lambda_{SM}/\sqrt{\xi}$ where Λ_{SM} is the unitarity violation scale in the SM without a Higgs. This implies that the Higgs mass can be raised to values approaching this number. In other words, the unitarity bound on the Higgs mass can be relaxed by about a factor $\sim \sqrt{1/\xi}$: if the Higgs only contributes 10% with respect to the SM, it can be as heavy as about 3 TeV; the bound is raised up to about 10 TeV if the contribution is only 1% ($\xi = 0.01$). We have plotted various contours for ξ in Fig. 1 (dashed lines).

It is well known that in Higgsless models particular care is needed to maintain compatibility with electroweak precision tests (EWPT) while simultaneously getting a large enough top quark mass. Concerning EWPT, the main problem turns out to be large corrections to the S parameter. In order to make them small, one has to allow the light fermions to be spread in the bulk [10]. When their profile is approximately flat, their wave function is orthogonal to the KK gauge bosons, and the contributions to the S parameter can be made arbitrarily small. The top mass problem has been successfully solved in ref. [9]. One has to use an alternative bulk-gauge/custodial representations for the third generation [14]. Instead of having the left-handed (LH) top and bottom transforming as a $(\mathbf{2}, \mathbf{1})$ under $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$ and the right-handed (RH) fields as a $(\mathbf{1}, \mathbf{2})$, one has to embed the LH top and bottom in a bidoublet $(\mathbf{2}, \mathbf{2})$, the RH top in a singlet $(\mathbf{1}, \mathbf{1})$ and the RH bottom in a triplet $(\mathbf{1}, \mathbf{3})$. Localizing the LH bidoublet and the RH top close to the TeV brane, while the RH bottom close to the Planck brane allows for produce a large enough top mass without generating large deviations to the $Zb_\ell\bar{b}_\ell$ coupling. More details are discussed in Appendix C. In the context of a bulk Higgs model exactly the same methods can be used to solve the two problems discussed above. Of course as $1/R'$ is increase to be larger than in the Higgsless limit, the bounds become less severe and thus a larger region of parameter space is allowed.

Another potential issue is if the theory is perturbative in every sector up to at least $\sim 5 - 10$ TeV, so that higher order corrections will not spoil EWPTs. We have to be concerned by the gauge, top and Higgs sector. Using the techniques of Naive Dimensional Analysis (NDA), properly adapted to 5D, one can estimate the scale at which each sector becomes strongly coupled. The Higgs sector, as long the lightest Higgs boson mode is lighter than about 1 TeV, becomes strongly coupled at a very high scale. However, the gauge and the top sector may still pose a potential problem. The strong scale of the weak gauge coupling was estimated in [10]: it crucially depends on the IR scale R'. Even taking into account factors as big as 4 [15], we see that the strong scale is safely above 10 TeV for $1/R' > 300$ GeV (a less conservative bound of 5 TeV on the strong scale would imply $1/R' > 160$ GeV, however this region is already unrealistic due to gauge boson resonances being too light). In

the top sector, we can estimate the strong scale corresponding to the bulk top Yukawa, y_5, using NDA (properly warped down):

$$\Lambda_{top} \sim \frac{24\pi^3}{y_5^2} \frac{R}{R'}\,, \tag{2.10}$$

where we are neglecting unknown factors of order 1. The situation, however, is more tricky, since the precise value of y_5 crucially depends on the localization of the left and right-handed tops. In Fig. 1 we show two thick gray lines corresponding to the top Yukawa becoming strong at 10 and 20 TeV, for $c_L = c_R = 0$ (this is in the range of parameters we use to produce a heavy top mass and small corrections to the $Zb_\ell \bar{b}_\ell$ coupling). It is interesting to notice that the more Higgsless the theory is, the lower the strong scale is. This is a consequence of a lower IR scale, and a smaller overlap with the top quark when we localize the Higgs on the brane. However, in the strict Higgsless limit $V \to \infty$ this scale becomes large again, due to the fact that in this limit the Yukawa coupling goes to zero. In this limit, the cut-off of the theory is set by the scale at which the gauge sector becomes strong. We stress again that the value of this scale strongly depends on the localization of the top quark. In the following we will be conservative, and consider only points where the non-perturbative scale is above 20 TeV.

3 Phenomenology

As in many other models with extra dimensions, bulk Higgs models predict the existence of KK states for all the SM fields. However, the most distinctive feature is the presence of a Higgs boson with an unconventional phenomenology. In the region where $V \simeq 246$ GeV the model approaches the SM, and the Higgs phenomenology is conventional. In the "Higgsless" limit ($V \gg 1/R', \beta \gg 1$) the Higgs decouples from the theory: even if we want to make it light, through a tiny quartic coupling (2.7), its couplings to other fields go to zero, and thus we will never be able to produce and observe it. It is interesting to study some intermediate points, where the couplings of the Higgs to the SM fields, although suppressed, are large enough to lead to a potential discovery at the LHC. In Fig. 1 we plotted the suppression of the coupling of the Higgs to gauge bosons. In Fig. 2 we show the suppression of the couplings with different SM fields as a function of V (for fixed $\beta = 2$). The first thing to notice is that the Higgs couplings are more suppressed to heavier particles. This can be understood in the following way. The heavier the particle, the more it couples to the VEV: this also implies that the wave functions are more distorted away from zero-modes (they are more repelled from the peak of the VEV). Schematically, the masses receive one contribution from the direct overlap with the VEV, and the other from terms involving the derivative along the extra dimension (this contribution is still induced by the VEV, as the light states would be zero modes for vanishing VEV). Thus, a larger part of their mass comes from the extra dimension compared to lighter particles. However, the (light) Higgs wave function is essentially identical to the VEV profile, so that the coupling is proportional to the portion of the mass coming directly from the VEV. As we already discussed, this portion is larger for
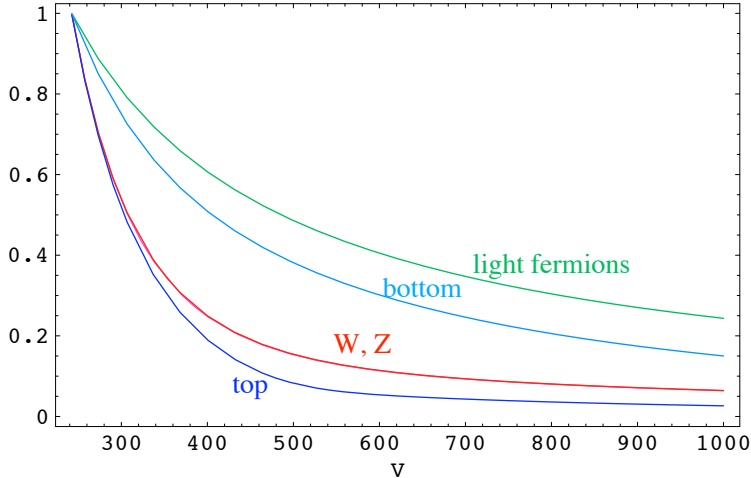
Figure 2: Suppression of the couplings of various SM fields with the Higgs with respect to the SM values for $\beta = 2$, as a function of V.

a) $V = 300$ GeV, $\beta = 2$	
$g_{ttH}/$SM	0.52
$g_{WWH}/$SM	0.54
$g_{ZZH}/$SM	0.54
$g_{bbH}/$SM	0.75
$g_{ffH}/$SM	0.81

b) $V = 500$ GeV, $\beta = 2$	
$g_{ttH}/$SM	0.08
$g_{WWH}/$SM	0.15
$g_{ZZH}/$SM	0.15
$g_{bbH}/$SM	0.38
$g_{ffH}/$SM	0.49

Table 1: Higgs couplings to the SM fields for the two benchmark points a) $V = 300$ GeV and $\beta = 2$, b) $V = 500$ GeV and $\beta = 2$. "f" stands for the light fermions

light fermions compared to heavy ones. We can therefore see that the couplings of the Higgs to heavy particles, like the top or weak gauge bosons, are more suppressed than to light fermions, like the bottom. This can affect the decay modes in an important way. Another interesting point is that the production of the Higgs (which mostly happens via the couplings to the W, Z and t) is approximatively suppressed by the "Higgslessness" of the model: the parameter ξ, plotted in Fig. 1. In other words, in a point where $\xi = 0.01$, the production rate of the Higgs is about 1% of the SM one.

To be more concrete, we choose two benchmark points: a) $V = 300$ GeV and $\beta = 2$, b) $V = 500$ GeV and $\beta = 2$. We report in Table 1 the couplings of the Higgs to various SM fields, compared to the SM, in the two cases, assuming a Higgs mass of 120 GeV (these numbers have a mild dependence on the Higgs mass below 1 TeV).

First of all, the suppressed HZZ coupling relaxes the LEP bound $m_H > 114$ GeV [16]. It turns out that for point a) the production cross section is suppressed by a factor around 4, and thus the bound becomes $m_H \gtrsim 95$ GeV [16]. For point b) the production rate is suppressed
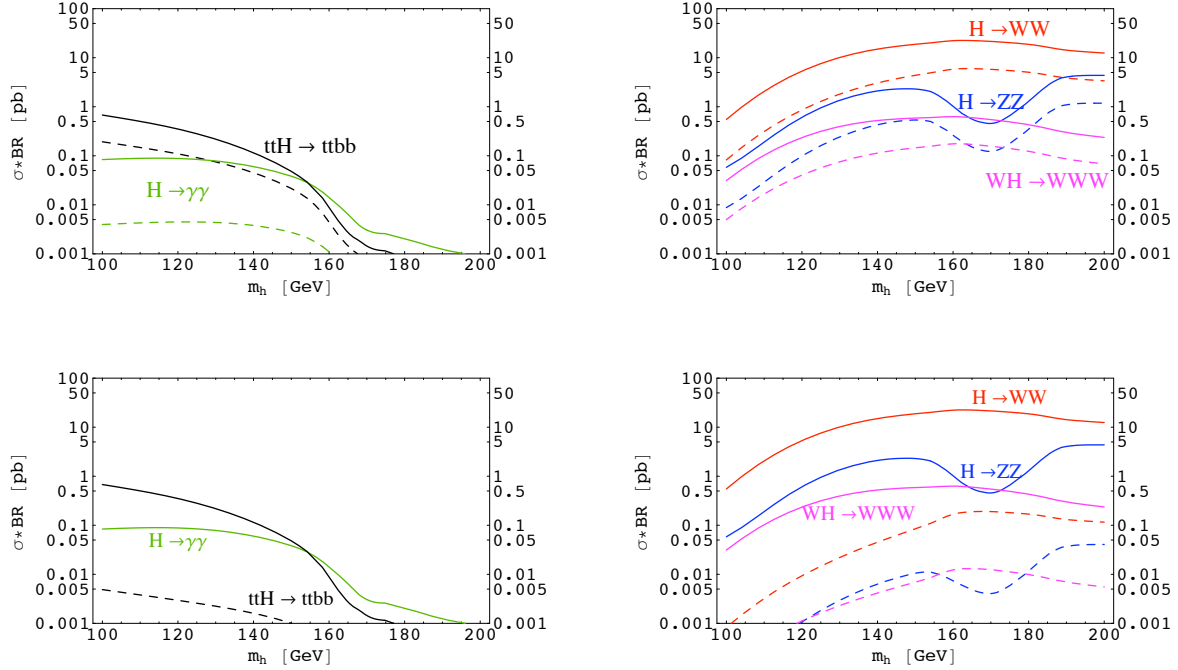
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Figure 3: Cross sections times branching ratios for various Higgs production and decay channels for the SM (solid lines) and gaugephobic Higgs (dashed lines) for $\beta = 2$ with $V = 300$ (top) and $V = 500$ (bottom).

by a factor of about 100, so there is no limit from LEP. At the LHC, the suppressed couplings to the top and the weak gauge bosons lead to a reduced production cross section. Both gluon fusion and weak boson fusion production channels receive a suppression of about a factor 4 in point a) and 100 in point b). Concerning the decay modes, branching fractions to light particles are enhanced with respect to the SM.

We show in Fig. 3 the production cross section times the branching fraction as a function of the Higgs mass for a single Higgs decaying to $WW^{(*)}, ZZ^{(*)}, \gamma\gamma$ and associated production $t\bar{t}H \to t\bar{t}b\bar{b}$, $WH \to WWW^{(*)}$. These are the most promising channels for the discovery at the LHC in various mass ranges. Continuous lines represent the SM, while dashed lines represent the gaugephobic Higgs model. It is important to notice the Higgs into two photons decay channel. It is mediated by a loop of tops and W's, which turn out to have the strongest suppression. For low Higgs masses (when the ZZ^* channel is still indistinguishable from $Z\gamma^*$ background) in the SM there are two possible discovery channels: $H \to \gamma\gamma$ and associated production $t\bar{t}H \to t\bar{t}b\bar{b}$, with a similar discovery significance. Since the coupling of the Higgs to t and W is more suppressed than to the b, the $H \to \gamma\gamma$ channel is no longer useful. All other channels get rescaled approximately by a factor 4 in point a) and 100 in point b). Since the discovery significance scales with the square root of the integrated luminosity, it means that at the LHC it will be necessary to collect 16 and 10^4 times the luminosity necessary

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a) $V = 300$ GeV, $\beta = 2$	
$1/R'$	372.5 GeV
W'	918 GeV
Z'_1	912 GeV
Z'_2	945 GeV
G'	945 GeV

b) $V = 500$ GeV, $\beta = 2$	
$1/R'$	244 GeV
W'	602 GeV
Z'_1	598 GeV
Z'_2	617 GeV
G'	617 GeV

Table 2: Spectrum of the first gauge boson KK states for the two benchmark points a) $V = 300$ GeV and $\beta = 2$, b) $V = 500$ GeV and $\beta = 2$.

in the SM case to claim a discovery with the same statistical significance. Let us discuss a concrete example. In the SM, discovering a 100 GeV Higgs in the $t\bar{t}H \to t\bar{t}b\bar{b}$ channel at a $\sim 8\sigma$ significance requires an integrated luminosity of about 100 fb^{-1} [17]. In this model, in case a), and with the same integrated luminosity, since the production rate decreases by a factor of 4 the significance would only be $\sim 2\sigma$. A 5σ discovery would require an integrated luminosity of $100 \cdot (5/2)^2 \sim 600$ fb^{-1}, which almost saturates the total integrated luminosity (700 fb^{-1}) that the LHC is expected to collect. The discovery might still be claimed, but the signal would be delayed by years and become statistically significant only at the very end of the LHC run. Point b) is much more problematic: the suppression factors are around 100, making the discovery of the Higgs at the LHC impossible.

As we discussed, there are regions in the parameter space where the LHC can discover the Higgs (even though requiring much more time than the SM Higgs), and regions where the Higgs is definitely out of reach for the LHC. It is therefore interesting, and complementary, to look for other new particles in the gauge sector: the gauge boson KK modes. As in the Higgsless model, they play a fundamental role in the gauge sector, restoring the perturbative unitarity in the longitudinal W and Z scattering. As a consequence, the "more Higgsless" the models is, the lighter those resonances are. In Table 2, we show the spectrum of the lightest modes for the two benchmark points. The phenomenology of these states was already studied in [18], where the authors only considered the couplings among gauge bosons. They showed that the signal should be observable for masses below ~ 1 TeV. Their approximation is valid as long as the couplings with fermions is small: this is true for light fermions [10], but not for the third generation of quarks. The reason is that top and bottom are localized towards the TeV brane, so that their wave functions have a larger overlap with the KK states (localized on the TeV brane as well). The study of vector resonances coupling negligibly to the light fermions but sizeably to the third generation of quarks has been performed in [19]. They show that, with an integrated luminosity of 300 fb^{-1}, a vector resonance can be discovered at 5σ significance up to about 2 TeV. The most promising channels are $gg \to b\bar{b}$, $t\bar{t}$, with a Z' radiated by one of the heavy quarks. The Z' would then decay mainly to top pairs, giving as a final state either $b\bar{b}t\bar{t}$ or $t\bar{t}t\bar{t}$. The choice of the couplings of the Z' to the third generation made in [19] is different from our case: they consider a Z' coupling only to right-handed top and bottom quarks. However the right handed couplings of the top and the bottom used in [19] approximately match the sum of left and right handed couplings of

charge	a) $V = 300$ GeV	b) $V = 500$ GeV	
5/3	581 GeV	382 GeV	X_L
2/3	643 GeV	511 GeV	T_L
-1/3	1062 GeV	712 GeV	b_R
2/3	1058 GeV	693 GeV	T_R
5/3	1124 GeV	832 GeV	X_R
2/3	1160 GeV	831 GeV	$t_L - t_R$
-1/3	1242 GeV	917 GeV	b_L
2/3	1318 GeV	1114 GeV	$t_L - t_R$

Table 3: Spectrum of the first resonances in the third generation quark sector, for the two benchmark points. Near the mass eigenvalues, we report the field where such eigenstate mostly live. The effect of the Yukawa couplings is indeed to mix the representations, but numerically it corresponds a small shift of the masses.

the third generation fermions in this model, so we expect their result to be approximately applicable to our model as well. Thus, for both our benchmark points the discovery of the Z' should not be missed at the LHC. The Z' can also be produced as a resonance in a Drell-Yan process. The coupling to the light quarks is suppressed, but is still non-zero. Compared to the previous case, one has one particle less in the final state, which can compensate the effect of the suppression. For this reason, the cross section might be comparable to the $Z't\bar{t}$ and $Z'b\bar{b}$ associated production. This channel would correspond to the search of a heavy resonance in the $t\bar{t}$ channel.

In the fermion sector, an interesting feature of this model is the structure of the resonances of the third generation quarks. This structure may be used to probe the new representations that minimize the deviations in the $Zb_\ell\bar{b}_\ell$ coupling. In table 3 we list the first fermionic resonances (below about 1 TeV) for the two benchmark points, and for a particular choice of the bulk masses ($c_L = c_R^t = 0$ and $c_R^b = -0.79$) that is compatible with current data. There are many particles that are light enough to be produced at LHC. The Yukawa interactions mix all the representations, however their effect is numerically small (order 10% in the masses) so that each particle is an approximate interaction eigenstate (this is more true for small V): this allows to try to reconstruct the non-trivial bulk gauge representations. For instance, for the benchmark point a), we can easily identify a doublet (X, T) with charges $+5/3$ and $+2/3$ respectively at ~ 600 GeV as the lightest particles. At ~ 1050 GeV we have a triplet of states that correspond to the triplet containing b_R, and around ~ 1200 GeV we have a doublet (t_L, b_L) and a singlet t_R. The discovery of a fermion with charge 5/3 alone would be a strong indication that a bidoublet exists, but from the spectra in the table we can see that a more detailed structure can actually be observed notwithstanding the mixings induced by the Yukawa couplings. The LHC should be able to distinguish the light doublet, and maybe more details of the spectrum depending on the precision in the mass determination. The heavy top will be pair produced in gluon fusion, or singly produced in Wb or Zt fusion ($gq \to T\bar{b}q'$ and $gq \to T\bar{t}q$). The techniques for the search of heavy partners

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of the top have been analyzed in the context of Little Higgs models. If the mass is below 1 TeV, as in our case, its discovery at the LHC cannot be missed [20]. The exotic quark X with charge $5/3$ is more interesting, as it is a direct consequence of the new realization of the custodial symmetry. Similarly to T, it can either be pair produced through gluon fusion or singly produced in the process $gq \to X\bar{t}q'$. The X would then decay to a tW^+ pair. The first mechanism would give 4 W's and 2 b-jets, while the second one 3 W's and 2 b-jets: for both processes the SM background would be very small, and thus we expect that the discovery at the LHC should be guaranteed. Due to its charge, in the decay chain there will be two same-sign W's, $X \to W^+t \to W^+W^+b$: in the leptonic channel they will lead to same-sign lepton pair events.

Another consequence of the new realization of the custodial symmetry in the bulk, is a suppression in the Wt_lb_l coupling, that can be parameterized in terms of a suppression of an effective CKM element V_{tb} [2]. If the assumption of the unitarity of the CKM matrix is relaxed, such suppression is not ruled out by present experiments [21]. The only way to directly probe this coupling is to measure single top production at the Tevatron and/or the LHC. In the two benchmark points, we find that the effective V_{tb} is 0.88 for $V = 300$ GeV and 0.74 for $V = 500$ GeV (to be compared with the SM value $V_{tb} \sim 1$). At the Tevatron, this implies a suppression of 23% in the benchmark point a) and 45% in point b) in the t-channel production cross section (and associated Wt production). In the s-channel, where the top is produced in association with a bottom quark via a virtual W, we should also take into account the interference with the W's: this might induce an even larger suppression [22]. At the LHC, the measurement of single top production in the t-channel will allow a direct measurement of V_{tb} at a 5% level [21], therefore testing directly this prediction of the model. The s-channel is more challenging to measure: due to the presence of 2 b-jets in the final state (one is from the top decay) it suffers from a large background from $t\bar{t}$ production.

Finally, the spectrum of KK excitations also contains massive gluons (the masses are listed in Table 2). It is not a very specific signature of this model, as it does not play any role in EWSB and it is only a consequence of putting the $SU(3)_C$ gauge group in the bulk. Its coupling to the third generation is enhanced for the same reason why the Z' coupling is: a significant overlap near the TeV brane. Thus it can be radiated away from a heavy quark. The signature is very similar to the Z' analyzed in [19] but with bigger couplings: now they are $SU(3)_C$ couplings rather than $SU(2)$. Thus the LHC is sensitive to even higher masses. In this case the discovery of the KK gluon G' at the LHC cannot be missed.

4 Conclusions

We have begun an exploration of the space of generalized RS scenarios where the Higgs localization width and VEV are taken as free parameters. Generically the Higgs is gaugephobic and topphobic. Such a gaugephobic Higgs can be lighter than the SM experimental bound or heavier than the SM theoretical bound. A gaugephobic Higgs will be more difficult (or

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[2]In the language of KK states, this reduced coupling is indeed the effect of the mixing of the top with heavy KK fermions via the Yukawa couplings.

impossible) to find at the LHC. However other particles (e.g. gauge boson resonances) are lighter than in conventional RS scenarios and thus will be easier to find. The Z' is a good example of a resonance that should be easy to look for. There are also top quark resonances and a tower of exotically charged fermions that is necessary for the implementation of the new realization of custodial symmetry. Since the top quark couplings are also reduced in gaugephobic Higgs scenarios, it is possible that the first experimental signature that can be found will be the suppression of single top production at the Tevatron.

Acknowledgements

A   The Higgs sector

The equations of motion for the Higgs field (with 4D momentum p_ν) and the BC's that arise from the variation of the action corresponding to Eq. 2.3 are:

$$\left(z^3 \partial_z \frac{1}{z^3} \partial_z + p^2 - \frac{\mu^2}{z^2} \right) \mathcal{H} = 0 \,, \tag{A.1}$$

$$\left(\frac{R}{R'} \right)^3 \partial_z \mathcal{H} + \frac{\partial}{\partial \mathcal{H}^*} V_{\text{TeV}} \bigg|_{R'} = 0 \,, \tag{A.2}$$

$$\partial_z \mathcal{H} - \frac{\partial}{\partial \mathcal{H}^*} V_{\text{UV}} \bigg|_{R} = 0 \,. \tag{A.3}$$

We will assume that the quartic term is localized on the IR brane, as we want EWSB to take place there.

In order to find the VEV we need to solve the equations of motion and find a solution that minimizes the full potential, as encoded in the BC's. For $p_\nu = 0$ and a diagonal VEV which breaks $\text{SU}(2)_L \times \text{SU}(2)_R \to \text{SU}(2)_D$:

$$< \mathcal{H} > = \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix} \frac{v(z)}{\sqrt{2}} \,. \tag{A.4}$$

the solutions are of the form

$$v(z) = a \left(\frac{z}{R} \right)^{2+\beta} + b \left(\frac{z}{R} \right)^{2-\beta} \,, \tag{A.5}$$

where $\beta = \sqrt{4 + \mu^2}$. Note that the dimension of the corresponding CFT operator is $2 \pm \beta$ depending on the UV boundary conditions. We can see that the effect of β, i.e. of the

bulk mass, is to control the localization of the profile of the VEV in the bulk. The bulk mass squared can be negative but it is bounded [23] by $\mu^2 > -4$. The form of the localized potentials $V_{\text{TeV}}, V_{\text{UV}}$ determines if a VEV is generated and its size. We want a VEV generated on the TeV brane, so we will add a "mexican hat" potential there. On the UV brane, on the other hand, we just add a mass term:

$$V_{\text{UV}} = m_{\text{UV}} \operatorname{Tr} |\mathcal{H}|^2 , \tag{A.6}$$

where m_{UV} has dimension of mass. Eq. A.3 can be used to fix the ratio between the two coefficients in the solution (A.5):

$$\frac{b}{a} = -\frac{2 + \beta - m_{\text{UV}} R}{2 - \beta - m_{\text{UV}} R} . \tag{A.7}$$

The relation between the mass m_{UV} and the bulk mass can select one of the two solutions: $m_{\text{UV}} R = 2 \pm \beta$ will select the solution growing towards the TeV or the UV brane respectively. In the CFT language, this boundary condition corresponds to the determining the CFT operator (and its scaling dimension) represented by the bulk Higgs. In the following we will assume that only the solution with the IR localized profile (corresponding to the operator with the higher dimension) is selected by the UV mass so that:

$$m_{\text{UV}} = \frac{2 + \beta}{R} \tag{A.8}$$

$$v(z) = a \left(\frac{z}{R}\right)^{2+\beta} . \tag{A.9}$$

Note that the dimension of the corresponding CFT operator which breaks electroweak symmetry is $2 + \beta$: being $\beta \geq 0$, in this case the Higgs mass is naturally of order TeV ($1/R'$). Tuning the UV mass as in Eq. A.8 is not essential: the solution we select is in fact growing faster towards the TeV brane, so it will have a larger effect on the W and Z mass. We will keep only one solution just to simplify our calculation. Note also that if we wanted to select the other solution ($z^{2-\beta}$), all we would have to do is to flip the sign of β in all the equations: in other words, taking $\beta < 0$ formally corresponds to the other solution. This will only be interesting for one reason: for $\beta = -1$ ($v \sim z$) we have a flat profile for the Higgs VEV.

On the TeV brane, we add the following potential:

$$V_{\text{TeV}} = \left(\frac{R}{R'}\right)^4 \frac{\lambda R^2}{2} \left(\operatorname{Tr} |\mathcal{H}|^2 - \frac{v_{\text{TeV}}^2}{2}\right)^2 , \tag{A.10}$$

where the warp factors have been added so that all the parameters have a natural scale set by R. Note also that a factor of R^2 has been added to make λ dimensionless, while v_{TeV} has dimension [mass]$^{3/2}$. Imposing the BC in Eq. (A.2), the coefficient a in Eq. (A.9) is fixed. We can thus write the profile of the VEV as

$$v(z) = \frac{1}{R^{3/2}} \left(R^3 v_{\text{TeV}}^2 - \frac{2(2+\beta)}{\lambda}\right)^{1/2} \left(\frac{z}{R'}\right)^{2+\beta} . \tag{A.11}$$

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The VEV $v(z)$ is real, so there is a solution only if v_{TeV} is large enough: a TeV brane localized negative squared mass has to be big enough to overcome the effect of the positive bulk mass. We can define a 4D VEV as the integral of the 5D VEV along the extra-dimension:

$$v_4^2 = \int_R^{R'} dz \left(\frac{R}{z} \right)^3 v^2(z). \tag{A.12}$$

In the limit where the VEV is localized on the IR branes, v_4 is the value of the 4D VEV. Thus the 5D VEV can be rewritten as:

$$v(z) = \sqrt{\frac{2(1+\beta)}{R^3(1-(R/R')^{2+2\beta})}} \, v_4 R' \left(\frac{z}{R'} \right)^{2+\beta} . \tag{A.13}$$

Notice that the natural size of v_4 is $\sim 1/R'$.

At this point we can parameterize the bulk Higgs through three free parameters: the bulk mass β, the 4D VEV v_4 and the quartic coupling λ. The first controls the localization of the profile, the second sets the size of the Higgs VEV, while the third controls the mass of the lightest physical Higgs mode (which we will come back to later).

The only parameter that carries a dimension of mass is v_4. In order to make its value more physically meaningful, we can rescale it, and reparameterize the Higgs VEV v as:

$$v(z) = \sqrt{\frac{2(1+\beta)\log R'/R}{(1-(R/R')^{2+2\beta})}} \, \frac{gV}{g_5} \frac{R'}{R} \left(\frac{z}{R'} \right)^{2+\beta} , \tag{A.14}$$

where g is the SM SU(2) gauge coupling, and V is the new input parameter we are defining. The reason for this choice is that in the flat VEV case ($\beta \to -1$) the contribution of the VEV in the equation of motion for the gauge bosons (B.7) is

$$\frac{R^2}{z^2} \frac{g_5^2}{4} v^2(z) \xrightarrow[\beta \to -1]{} \frac{g^2 V^2}{4} ; \tag{A.15}$$

Since for $\beta \to -1$ the W wave function is flat, the W mass comes entirely from the Higgs VEV. Thus in this limit V has to be equal to the SM VEV $v_{SM} \sim 246$ GeV. We numerically checked that we obtain the SM when $V \to v_{SM}$ independently of the localization of the VEV.

The bulk Lagrangian for the physical Higgs h is:

$$\mathcal{L}_h = \int_R^{R'} dz \left(\frac{R}{z} \right)^3 \left\{ \frac{1}{2} (\partial_\mu h)^2 - \frac{1}{2} (\partial_z h + \partial_z v)^2 - \frac{1}{2} \frac{\mu^2}{z^2} (h+v)^2 \right\} . \tag{A.16}$$

The tadpoles in h cancel out due to the equations of motion of the Higgs VEV $v(z)$. On the other hand, the localized potentials do generate a mass term for the Higgs. The equations of motion for a mode with $p_\nu p^\nu = m^2$ are:

$$\left(z^3 \partial_z \frac{1}{z^3} \partial_z + m^2 - \frac{\mu^2}{z^2} \right) h = 0 ,$$
$$\partial_z h - m_{\text{UV}} \, h|_R = 0 , \tag{A.17}$$
$$\partial_z h + (R/R') \, m_{\text{TeV}} \, h|_{R'} = 0 .$$

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The mass on the UV brane is given by Eq. (A.8). The solution of the bulk equation of motion and the UV brane boundary condition is given by

$$h_m(z) = A z^2 \left(Y_{1+\beta}(mR) J_\beta(mz) + J_{1+\beta}(mR) Y_\beta(mz) \right) , \tag{A.18}$$

where A is a normalization factor. The mass eigenvalues m are determined by the TeV localized potential. On the TeV brane, the effective mass term is related to the Higgs VEV and the quartic coupling λ:

$$m_{\text{TeV}} R = \lambda R^3 v^2(R') - (2+\beta) . \tag{A.19}$$

The value of this parameter determines the Higgs spectrum. As in the SM, the Higgs mass is a free parameter, determined by λ. However, λ will also determine the Higgs quartic coupling, so we need to make sure it is not too large in order to avoid strong coupling in the Higgs sector.

B Gauge sector

We now analyze the gauge sector of this model. Expanding the bulk scalar around the VEV we have:

$$\mathcal{H} = \frac{1}{\sqrt{2}} \left(v(z) + h \right) \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix} e^{i\pi^a \sigma^a} . \tag{B.1}$$

The π^as are the Goldstone bosons eaten by the broken gauge directions, associated with the combination of fields $A_L^a - A_R^a$. We will focus on those gauge fields and Goldstones: to keep the formulae simple and clear, we will use the notation of a simple group completely broken by the VEV, and a generalization to our model will be straightforward. The Lagrangian up to quartic terms becomes:

$$\mathcal{L}_{gauge} = \int_R^{R'} dz \, \frac{R}{z} \frac{1}{g_5^2} \left\{ -\frac{1}{4} F_{\mu\nu}^2 + \frac{1}{2} (\partial_z A_\mu)^2 + \frac{1}{2} \left(\frac{R}{z} \right)^2 g_5^2 v^2 A_\mu^2 + \right.$$

$$+ \frac{1}{2} (\partial_\mu A_5)^2 + \frac{1}{2} \left(\frac{R}{z} \right)^2 g_5^2 v^2 (\partial_\mu \pi)^2 - \frac{1}{2} \left(\frac{R}{z} \right)^2 g_5^2 v^2 (\partial_z \pi + A_5)^2 +$$

$$\left. - \left(z \partial_z \frac{A_5}{z} + \left(\frac{R}{z} \right)^2 g_5^2 v^2 \pi - A_5 \delta(z-R) + A_5 \delta(z-R') \right) \partial_\mu A^\mu \right\} , \tag{B.2}$$

where g_5 is the 5D gauge coupling of the SU(2)'s. Note that the Higgs VEV always appears in the following combination:

$$\tilde{v}(z) = \frac{R}{z} g_5 v(z) = \sqrt{\frac{2(1+\beta)}{(1-(R/R')^{2+2\beta})}} \frac{g_5}{\sqrt{R}} v_4 \left(\frac{z}{R'} \right)^{1+\beta} . \tag{B.3}$$

This quantity carries more physical meaning, as it is equivalent to a bulk mass for the gauge bosons. In general the gauge boson masses have two sources: the curvature of their wave function and the Higgs VEV. Notice also that if $\beta = -1$ the VEV is flat, the gauge boson wave function is also flat, and the mass term becomes constant:

$$\tilde{v}(z) \rightarrow \frac{g_5}{\sqrt{R \log R'/R}} v_4 = g_4 v_4 \,. \tag{B.4}$$

In this case the only source for the gauge boson mass is the Higgs VEV.

The mixing terms on the third line of Eq. B.2 can be cancelled out by R_ξ gauge fixing terms in the bulk and on the branes:

$$\mathcal{L}_{GF} = \int_R^{R'} dz \, -\frac{R}{z} \frac{1}{g_5^2} \left\{ \frac{1}{2\xi} \left(\partial_\mu A^\mu - \xi \left(z\partial_z \frac{A_5}{z} + \tilde{v}^2 \pi \right) \right)^2 + \right.$$

$$\left. +\frac{1}{2\xi_{\text{UV}}} (\partial_\mu A^\mu - \xi_{\text{UV}} A_5)^2 \, \delta(z - R) + \frac{1}{2\xi_{\text{TeV}}} (\partial_\mu A^\mu + \xi_{\text{TeV}} A_5)^2 \, \delta(z - R') \right\} \,. \tag{B.5}$$

The unitary gauge is, as usual, given in the limit where all the three ξ-parameters are sent to infinity. A combination of the π's and A_5's is eaten by massive gauge bosons, while another one gives rise to a tower of physical states. However, those scalars are as heavy as the KK states, and the do not play any interesting role in the symmetry breaking. In the unitary gauge, the vector bosons obey the following equation of motion:

$$\left(z\partial_z \frac{1}{z}\partial + m^2 - \tilde{v}^2(z) \right) A_\mu(z) = 0. \tag{B.6}$$

Ii is now straightforward to write the equations of motion for the gauge fields in our specific model:

$$\left(z\partial_z \frac{1}{z}\partial + m^2 - \frac{g_5^2}{4} \frac{R^2}{z^2} v^2 \right) (A_{L\mu}^a - A_{R\mu}^a) = 0 \,, \tag{B.7}$$

$$\left(z\partial_z \frac{1}{z}\partial + m^2 \right) (A_{L\mu}^a + A_{R\mu}^a) = 0 \,, \tag{B.8}$$

$$\left(z\partial_z \frac{1}{z}\partial + m^2 \right) B_\mu = 0 \,. \tag{B.9}$$

As in the Higgsless model, the BC's on the UV brane break $\text{SU}(2)_R \times \text{U}(1)_X \rightarrow \text{U}(1)_Y$:

$$\begin{aligned} \partial_5 A_{L\mu}^a &= 0 \,, & A_{R\mu}^\pm &= 0 \,, \\ \partial_5 \left(\frac{1}{g_5^2} B_\mu + \frac{1}{g_5^2} A_{R\mu}^3 \right) &= 0 \,, & A_{R\mu}^3 - B_\mu &= 0 \,. \end{aligned} \tag{B.10}$$

On the TeV brane all the gauge fields have Neumann BC's since EWSB is accomplished by the Higgs VEV.

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C The fermion sector

The conventional choice for embedding fermions into $SU(2)_L \times SU(2)_R$ models is to put the LH and RH SM fermions into $SU(2)_L$ and $SU(2)_R$ bulk doublets respectively. This is what we will be using for the light fermions. For instance, $\Psi_L = (u_L, d_L)$, $\Psi_R = (u_R, d_R)$ transforming as

$$(\mathbf{2}, \mathbf{1})_{1/6}, \quad (\mathbf{1}, \mathbf{2})_{1/6} \tag{C.1}$$

of $SU(2)_L \times SU(2)_R \times U(1)_X$. For these representations, the X-charge can be identified as $X = (B - L)/2$. The SM zero modes [3, 4, 24] can be reproduced by the assignment of the following BC's:

	Ψ_L	UV	IR		Ψ_R	UV	IR
$\chi_L = \begin{pmatrix} \chi_{u_L} \\ \chi_{d_L} \end{pmatrix}$		$+$	$+$	$\chi_R = \begin{pmatrix} \chi_{u_R} \\ \chi_{d_R} \end{pmatrix}$		$-$	$-$
$\psi_L = \begin{pmatrix} \psi_{u_L} \\ \psi_{d_L} \end{pmatrix}$		$-$	$-$	$\psi_R = \begin{pmatrix} \psi_{u_R} \\ \psi_{d_R} \end{pmatrix}$		$+$	$+$

$$\tag{C.2}$$

where $+$ stands for a Neumann BC, $-$ stands for a Dirichlet BC, χ represents the LH chirality, and ψ the RH chirality. To give the fermions a mass, a bulk Yukawa coupling can be written. After replacing the Higgs with its VEV the mass term has the form

$$\frac{\lambda}{\sqrt{2}} v(z) \left(\chi_L \psi_R + \chi_R \psi_L \right) + h.c. \tag{C.3}$$

. Note that in order to preserve custodial symmetry, the Yukawa coupling (C.3) couples to both up and down type quarks. To split their mass, a large kinetic term can be added on the UV brane for ψ_{d_R}, since $SU(2)_R$ is broken on that brane.

For the third generation, things are more problematic due to the large top mass. In order to generate the top mass, one needs to localize it closer to the IR brane, where the Higgs lives. However this turns out to generate large corrections to the $Zb_\ell \bar{b}_\ell$ coupling. To solve this problem, one has to use a different set of representations for the third generation, as proposed by [14]. The LH doublet is now embedded in a bidoublet of $SU(2)_L \times SU(2)_R$, the RH top in a singlet, and the RH bottom in a triplet of $SU(2)_R$:

$$\Psi_L = (\mathbf{2}, \mathbf{2})_{2/3} = \begin{pmatrix} t_L & X_L \\ b_L & T_L \end{pmatrix}, \qquad \Psi_R = (\mathbf{1}, \mathbf{3})_{2/3} = \begin{pmatrix} X_R \\ T_R \\ b_R \end{pmatrix}, \qquad t_R = (\mathbf{1}, \mathbf{1})_{2/3}, \tag{C.4}$$

where all these fermion fields are bulk fields. As usual, $Y = T_R^3 + X$, $Q = T_L^3 + Y$, and the BC's ensure that the only zero modes correspond to SM fields. There are two $SU(2)_L \times SU(2)_R \times U(1)$ invariant Yukawa couplings that one can write in the bulk. After plugging in the VEV for the bidoublet Higgs field these will lead to the bulk mass terms for the fermions

$$\mathcal{L}_Y = \lambda_3 \frac{v(z)}{\sqrt{2}} \left[\frac{1}{\sqrt{2}} T_R \left(t_L + T_L \right) + b_R b_L + X_R X_L \right] + \frac{\lambda_1 v(z)}{2} t_R \left(t_L - T_L \right) + h.c. \tag{C.5}$$

Under the unbroken $SU(2)_D$ subgroup, the combination $(t_L - T_L)/\sqrt{2}$ is the singlet component of Ψ_L, while the fields $(X_L, (t_L + T_L)/\sqrt{2}, b_L)$ form the triplet. Choosing Ψ_L and t_R to be localized close to the IR brane, the top mass can easily obtained. In order to avoid large deviations to the $Zb_\ell\bar{b}_\ell$, one has also to localize Ψ_R close to the UV brane [9]. Changing V and β changes the allowed parameter space, which becomes larger for larger values of $1/R'$.

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