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# Gravity and electroweak symmetry breaking in a RSI/RSII hybrid model

Brian Glover* and Jong Anly Tan†

*Particle Theory Group, Department of Physics, College of William and Mary, Williamsburg, Virginia 23187-8795, USA*

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We present a hybrid RSI/RSII model in which we both solve the hierarchy problem and produce a continuum of Kaluza-Klein graviton modes. In this model, four-dimensional gravity can be reproduced, and the radion mode can be stabilized. We then modify the hybrid gravity model to include $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ bulk gauge fields. Electroweak symmetry is broken by the choice of appropriate boundary conditions. By adjusting the size of one region of the extra dimension, we show that the $S$ parameter can be decreased while protecting the $\rho$ parameter from corrections. We find that as the $S$ parameter is decreased by ~60%, $M_{Z'}$ and $M_{W'}$ stay below 1800 GeV, protecting unitarity.

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## I. INTRODUCTION

The idea of warped extra dimensions was first introduced in 1983 when Rubakov and Shaposhnikov suggested that a vanishing four-dimensional cosmological constant would result if a five-dimensional bulk vacuum energy was tuned to cancel the large four-dimensional vacuum energy of the standard model (SM) fields [1]. This work was popularized in 1999 when Randall and Sundrum introduced two famous examples of warped extra dimensions, which led to interesting and distinct phenomenology (hereafter called RSI [2] and RSII [3]). In the first model (RSI), a finite warped extra dimension living between a positive and a negative tension brane was used to solve the hierarchy problem. This model predicts Kaluza-Klein (KK) graviton excitations to have masses on the order of a few TeV which could possibly be detected at the Large Hadron Collider (LHC) in the near future. In the RSII model, Randall and Sundrum considered a warped infinite extra dimension. Although they no longer solved the hierarchy problem, they found that four-dimensional gravity can still be reproduced in an infinite extra dimension since the corrections to Newton's law at large distances are suppressed on the positive tension brane.

Since these models were first introduced, many extensions of their work have been proposed. Some of these extensions include adding extra branes to the bulk of RSII [4–6], localizing gravity on thick branes [7], adding SM fields to the bulk of RSI [8], Higgless models in an RSI background [9], etc. In one of these models [5], an extra negative tension brane was included in the bulk of the infinite extra dimension of RSII. This model, if stable, was designed to solve the hierarchy problem as in RSI but with an infinite extra dimension. However, it was found that when the scalar gravity mode (radion) of the five-dimensional graviton is carefully considered, the theory becomes unstable [10]. This instability arose since the kinetic term of the radion in these theories was found to be negative [11]. The bulk stress tensor violates the positivity of energy condition and the brane is unstable to crumpling. More recently, Agashe *et al.* [12] pointed out that if one could stabilize a IR-UV-IR model with $Z_2$ parity about the UV brane, one could address the hierarchy problem naturally. They argue that in an alternate UV-IR-UV model, one would have to add large brane kinetic terms in order to solve the hierarchy problem. In Sec. II we propose a model in which the negative tension brane is placed at an orbifold fixed point with positive tension branes living in the bulk of an infinite, warped extra dimension (see Fig. 1). The metric is given by $ds^2 = e^{-A(y)}dx^2 + dy^2$ where the warp factor is

$$A(y) = \begin{cases} -2k_1|y| & \text{if } 0 \le |y| \le r \\ 2k_2|y| - 2(k_1 + k_2)r & \text{if } |y| > r. \end{cases} \quad (1.1)$$

As in Lykken and Randall [4], this theory has a continuous KK spectrum, while also solving the hierarchy problem. However, the phenomenology of our model is more of a hybrid between RSI and RSII in which the KK gravitons of RSI become resonances. Placing a negative tension brane

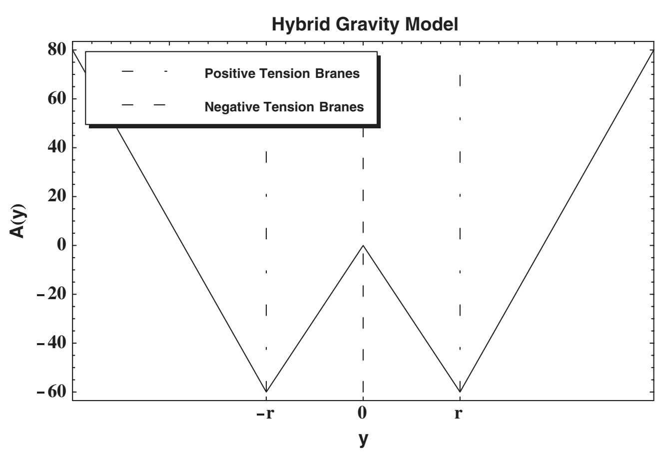


FIG. 1. The hybrid RSI/RSII gravity model. The space is orbifolded around $y = 0$ and extends to infinity.

*baglov@wm.edu
†jmtanx@wm.edu

 

at an orbifold fixed point projects out the negative energy mode of the radion and therefore allows the theory to be stabilized. We calculate the gravitational spectrum and show how this theory can be stabilized.

Warped extra dimensions have also proven to be interesting for models of Higgsless electroweak symmetry breaking. In Cacciapaglia *et al.* [9], $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge fields were included in the bulk of anti-de Sitter space (AdS) space. Custodial isospin was preserved by breaking $SU(2)_L \times SU(2)_R$ down to $SU(2)_D$ on the negative tension brane [13], while $SU(2)_R \times U(1)_{B-L}$ was broken down to $U(1)_Y$ on the Planck brane. It was found, as in technicolor theories, that an order one $S$ parameter is produced in conflict with experiments. In order to address this problem, a Planck brane kinetic term was added, which was found to decrease the $S$ parameter but at the price of destroying unitarity. They also added a $U(1)_{B-L}$ brane kinetic term to the TeV brane, which also lowered the $S$ parameter but at the price of making $T$ nonzero. More recently, Carone *et al.* [14] showed that a holographic UV-IR-UV model can be constructed, with $SU(2)_L \times U(1)_{B-L}$ gauge fields in the bulk, in which a custodial symmetry is generated without introducing a $SU(2)_R$ gauge group. They found that like the standard Higgsless model, the $S$ parameter is too large. In Sec. III, we modify our hybrid model to include gauge fields in the warped extra dimension. Following Csàki *et al.* [9], we include $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge fields in the bulk and use boundary conditions to break the symmetry in order to reproduce the SM on one of our branes. In order to have a normalizable photon, we have brought in another negative tension brane from infinity to cut off the space at an orbifold fixed point (see Fig. 2). We find corrections to the $\rho$ parameter to be suppressed, signaling that an approximate custodial symmetry is preserved. We calculate oblique corrections in this model and find that as the added slice of the extra dimension increases, the $S$ parameter decreases. We stress that this method of reducing the $S$ parameter appears to keep corrections to the $\rho$ parameter suppressed, while preserving unitarity for a decrease in $S$ up to 60%.

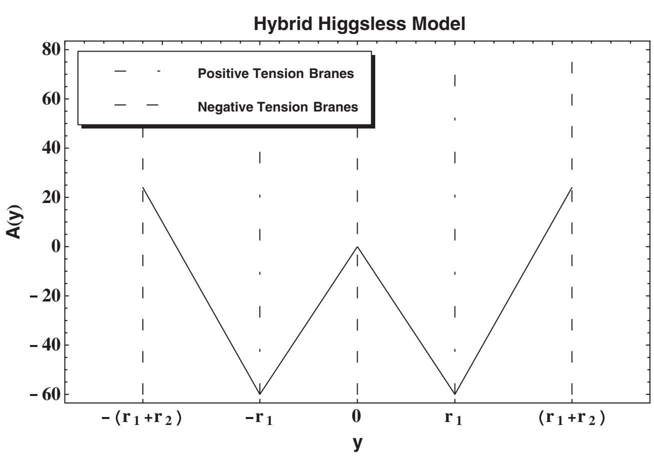


FIG. 2. The hybrid RSI/RSII Higgsless model. The space is orbifolded around $y = 0$ and ends at the location of the outside negative tension branes ($y = \pm(r_1 + r_2)$).

## II. GRAVITY IN THE HYBRID MODEL

Our theory is defined by placing a negative tension brane at an orbifold fixed point ($y = 0$) in an infinite fifth dimension (the TeV brane). Two additional positive tension branes are added at the points $y = \pm r$ (the Planck branes). It is important to point out that unlike the theories proposed in [4,5], we place the TeV brane at the orbifold fixed point, which (as we will discuss later) stabilizes the Radion mode (see [11]). The $Z_2$ symmetry demands that the tensions of the two additional Planck branes be equal. The action takes the form

$$S = \int d^5x\sqrt{-g^{(5)}}[2M_{\rm pl}^{(5)3}R - \Lambda_b - \sqrt{-g^{(4)}}V_1\delta(y) - \sqrt{-g^{(4)}}V_2\{\delta(y+r) + \delta(y-r)\}]. \qquad (2.1)$$

If we assume four-dimensional Poincaré invariance, the metric is given by

$$ds^2 = g_{MN}dx^M dx^N, \qquad (2.2)$$

with

$$g_{MN}(x^\mu, y) = \begin{pmatrix} -e^{-A(y)} & 0 & 0 & 0 & 0 \\ 0 & e^{-A(y)} & 0 & 0 & 0 \\ 0 & 0 & e^{-A(y)} & & 0 \\ 0 & 0 & 0 & e^{-A(y)} & 0 \\ 0 & 0 & 0 & 0 & 1 \end{pmatrix} \qquad (2.3)$$

and

$$A(y) = \begin{cases} -2k_1|y| & \text{if } 0 \le |y| \le r \\ 2k_2|y| - 2(k_1 + k_2)r & \text{if } |y| > r. \end{cases} \qquad (2.4)$$

As in [2], the assumption of four-dimensional Poincaré invariance leads one to derive the tension of the TeV brane located at $y=0$ to be $V_1 = -24M_{\rm pl}^{(5)3}k_1$, and the cosmological constant between the Planck and TeV branes is $\Lambda_1 = -24M_{\rm pl}^{(5)}k_1^2$. Likewise, the tension on the Planck brane located at $y=r$ is found to be $V_2 = 24M_{\rm pl}^{(5)3} \times (k_1 + k_2)$, and the cosmological constant outside the Planck branes is $\Lambda_2 = -24M_{\rm pl}^{(5)}k_2^2$. It is useful to transform the metric to manifestly conformally flat coordinates, where Einstein's equations take a simpler form. In these coordinates, the metric takes the form

$$g_{MN}(x^\mu, z) = e^{-A(z)}\mathrm{diag}(-1, 1, 1, 1, 1), \qquad (2.5)$$

where

$$e^{-A(z)} = \begin{cases} \frac{1}{(-k_1|z|+1)^2} & \text{if } z \leq z_b \\ \frac{1}{(k_2|z|+C)^2} & \text{if } z > z_b. \end{cases} \quad (2.6)$$

Now the Planck branes are located at $z_b = \pm(1 - e^{-k_1/r})/k_1$, and the constant $C = -k_2/k_1 + \exp[-k_1 r](1 + k_2/k_1)$ is chosen such that $z_b$ is the same for the two slices of AdS space.

### A. Kaluza-Klein modes

For now we will just consider the spin-2 fluctuation of the metric. The scalar mode (radion) will be discussed in the following section. Consider a perturbation of the form

$$ds^2 = e^{-A(z)}(dx^\mu dx^\nu(\eta_{\mu\nu} + h_{\mu\nu}(x,z)) + dz^2). \quad (2.7)$$

The transverse traceless solution can be written as $h_{\mu\nu}(x,z) = e^{3A(z)/4}\tilde{h}_{\mu\nu}(x)\psi(z)$, where $\Box_4 \tilde{h}_{\mu\nu}(x) = m^2\tilde{h}_{\mu\nu}(x)$ and

$$[-\partial_z^2 + V(z)]\psi(z) = m^2\psi(z). \quad (2.8)$$

The potential $V(z)$ is found to be [3]

$$\begin{aligned} V(z) &= \frac{9}{16}(\partial_z A(z))^2 - \frac{3}{4}\partial_z^2 A(z) \\ &= \{\begin{matrix} \frac{15k_1^2}{4(-k_1|z|+1)^2} & \text{if } |z| \leq zb \\ \frac{15k_2^2}{4(k_2|z|+C)^2} & \text{if } |z| > zb \end{matrix}\} + \frac{3k_1}{(-k_1+1)}\delta(z) \\ &\quad - \frac{3}{2}\left(\frac{k_1}{(-k_1|z|+1)} + \frac{k_2}{(k_2|z|+C)}\right) \\ &\quad \times (\delta(z - z_b) + \delta(z + z_b)). \end{aligned}$$

As usual, since the equation of motion for the Kaluza-Klein modes can be written in the form $\hat{Q}^\dagger\hat{Q}\psi(z) = m^2\psi(z)$, with $\hat{Q} = \partial_z + (3/4)A'(z)$, there is a zero mode

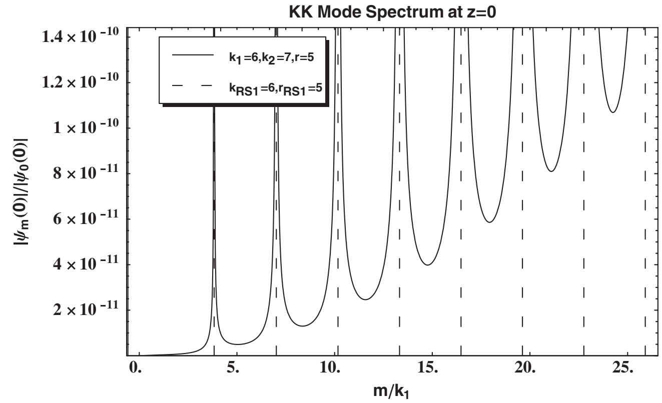


FIG. 3. Mass spectrum for both the hybrid RS (solid) and RSI (dashed) models. The hybrid RS model's spectrum was normalized by the zero mode's value at $z = 0$.

solution that satisfies $\hat{Q}\psi_0(z) = 0$:

$$\psi_0(z) = N\exp[-\tfrac{3}{4}A(z)]. \quad (2.9)$$

$N$ is found by normalization $N = [\int \exp[-3/2A(z)]dz]^{-1/2}$.

The higher KK modes are found by solving Eq. (2.8) subject to the following boundary conditions and normalization:

(1) $\psi_m(z)$ is continuous at the Planck branes ($z = \pm z_b$).
(2) $\psi'_m(z)$ is discontinuous at
(a) the TeV brane: $\Delta(\psi'_m(z))|_{z=0} = 3k_1\psi_m(0)$.
(b) the Planck branes $\Delta(\psi'_m(z))|_{z=\pm z_b} = -3/2(\frac{k_1}{-k_1|z_b|+1} + \frac{k_2}{k_2|z_b|+C})\psi_m(\pm z_b)$.
(3) $\psi_m(z)$ approaches a normalized plane wave solution for very large $z$.

The solution is

$$\psi_m(z) = \begin{cases} (-|z| + 1/k_1)^{1/2}[a_m Y_2(m(-|z| + 1/k_1)) + b_m J_2(m(-|z| + 1/k_1))] & \text{if } |z| \leq zb \\ (|z| + C/k_2)^{1/2}[a'_m Y_2(m(|z| + C/k_2)) + b'_m J_2(m(|z| + C/k_2))] & \text{if } |z| > zb. \end{cases} \quad (2.10)$$

The boundary conditions and normalization give the following relationships among the coefficients:

$$\begin{aligned} &a_m Y_2(me^{-k_1 r}/k_1) + b_m J_2(me^{-k_1 r}/k_1) \\ &= \left(\frac{k_1}{k_2}\right)^{1/2}[a'_m Y_2(me^{-k_1 r}/k_2) + b'_m J_2(me^{-k_1 r}/k_2)], \end{aligned} \quad (2.11)$$

$$a_m Y_1(m/k_1) + b_m J_1(m/k_1) = 0, \quad (2.12)$$

$$\begin{aligned} &a_m Y_1(me^{-k_1 r}/k_1) + b_m J_1(me^{-k_1 r}/k_1) \\ &= \left(\frac{k_1}{k_2}\right)^{1/2}[a'_m Y_1(me^{-k_1 r}/k_2) + b'_m J_1(me^{-k_1 r}/k_2)], \end{aligned} \quad (2.13)$$

$$a'^2_m + b'^2_m = m. \quad (2.14)$$

Unlike the RSI model, there is a continuous spectrum of graviton modes (all $m > 0$ are allowed). The RSI spectrum is discrete and given by $m_n = k_1 x_n$, where $x_n$ denotes the zeros of $J_1(x)$ [15].[1] In Fig. 3 we compare the hybrid RS KK spectrum to that of RSI. We have chosen order one parameters such that $k_1 r = 30$. The resonances in the spectrum correspond nicely to the discrete spectrum found in RSI. Since the modes are suppressed compared to the zero mode, the corrections to Newton's law are small:

[1] Since we have normalized the metric to be 1 at the TeV brane instead of the Planck brane as done in RSI [2], our spectrum is multiplied by $\exp[k_1 r]$ as compared to the solution found in [15].

$$V(\bar{x}, z = 0, \bar{x}', z' = 0) = \frac{1}{2M_{\rm pl}^3} \frac{|\psi_0(0)|^2}{|\bar{x} - \bar{x}'|} + \int_0^\infty \frac{1}{2M_{\rm pl}^3} \times \frac{|\psi_m(0)|^2 e^{-m|\bar{x}-\bar{x}'|}}{|\bar{x} - \bar{x}'|} dm, \tag{2.15}$$

$$\sim \frac{1}{2M_{\rm pl}^3} \frac{|\psi_0(0)|^2}{|\bar{x} - \bar{x}'|}\left(1 + \int_0^\infty e^{-m|\bar{x}-\bar{x}'|} \frac{|\psi_m(0)|^2}{|\psi_0(0)|^2} dm\right). \tag{2.16}$$

### B. Radion stabilization

As mentioned above, placing the TeV brane at the orbifold fixed point will allow the radion mode to be stabilized. To see this we need to include the spin-0 fluctuation ($f(x)$) of the five-dimensional graviton. The proper way to include this mode was discussed in [11,16]. Starting in a gauge where both branes are flat, the metric can be parameterized as [11]

$$ds^2 = a(y, x)^2[\eta_{\mu\nu} + h_{\mu\nu} + 2\epsilon(y)\partial_\mu\partial_\nu f(x)]dx^\mu dx^\nu + b(y, x)^2 dy^2; \tag{2.17}$$

$$a(y, x) = e^{-A(y)}[1 + B(y)f(x)], \tag{2.18}$$

$$b(y, x) = \frac{(\partial_y \log a)^2}{A'(y)^2/4}, \tag{2.19}$$

where $\epsilon(y)$ depends on the coordinate choice and has $\partial_y\epsilon$ fixed at the position of the branes by matching conditions. It was found that with the additional relation

$$B(y) = 2e^{A(y)} + e^{-A(y)}A'(y)\epsilon'(y) - e^{A(y)}\int_0^y e^{-2A(\tilde{y})}A''(\tilde{y})\epsilon'(\tilde{y})d\tilde{y}, \tag{2.20}$$

the spin-2 calculation goes through as done in the previous section and is decoupled from the spin-0 radion mode ($f(x)$). For the warp factor given in Eq. (2.4), Pilo *et al.* found that the effective four-dimensional Lagrangian contains the term [11]

$$\mathcal{L} \supset 2M_{\rm pl}^{(5)3}\int_0^r dy \frac{3B'(y)}{A'(y)/2}(f\Box f), \tag{2.21}$$

$$= \frac{24M_{\rm pl}^{(5)3}}{2k_1}\left(1 - e^{-2k_1 r}\frac{k_2}{k_2 + k_1}\right)\int \sqrt{-g}d^4x f\Box f, \tag{2.22}$$

where the values of $B(y)$ are fixed at the brane positions by matching conditions. Notice that for positive definite $k_1$ and $k_2$, the kinetic term is always positive in our model. This should be compared to models with a negative tension brane in the bulk [4,16]. For these models, $k_1$ is opposite in sign, and the kinetic term therefore can become negative. Models with a negative tension brane in the bulk may therefore contain a ghost radion mode. In our model, however, the radion mode can be stabilized using a mechanism like the one introduced by Goldberger and Wise [17]. In addition, since our model has a normalizable graviton zero mode and lacks a ghost radion mode, we do not expect IR modifications to gravity [10,11].

## III. HIGGSLESS SYMMETRY BREAKING IN THE HYBRID MODEL

In this section we will put $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge fields in the bulk. The metric is given by (2.5) [see Fig. 2]. However, unlike before, in this section we cut off the infinite extra dimension in order to make the massless mode normalizable.[2] This is accomplished by adding a negative tension brane at an orbifold fixed point $y = (r_1 + r_2)$ [or $z = z_{b2} = 1/k_2(e^{(k_2r_2-k_1r_1)} - k_2/k_1(e^{-k_1r_1} - 1 + k_1/k_2 e^{-k_1r_1}))$ in $z$ coordinates]. The five-dimensional action for this model is

$$S = \int d^4x \int dz\sqrt{-g^{(5)}}\left[-\frac{1}{4}R^a_{MN}R^{aMN} - \frac{1}{4}L^a_{MN}L^{aMN} - \frac{1}{4}B_{MN}B^{MN}\right], \tag{3.1}$$

where $R^a_{MN}$, $L^a_{MN}$, and $B_{MN}$ are the $SU(2)_L$, $SU(2)_R$, and $U(1)_{B-L}$ field strengths.

Using the same procedure as [9], we choose to work in the unitary gauge where all KK modes of the fields $L^a_5$, $R^a_5$, $B_5$ are unphysical. Boundary conditions were imposed to break the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ symmetry to the standard model at $z = z_{b2}$ and to $SU(2)_D \times U(1)_{B-L}$ at $z = 0$. The boundary conditions are

$$z = 0: \begin{cases} \partial_z(L^a_\mu + R^a_\mu) = 0, L^a_\mu - R^a_\mu = 0, \partial_z B_\mu = 0, \\ L^a_5 + R^a_5 = 0, \partial_z(L^a_5 - R^a_5) = 0, B_5 = 0 \end{cases} \tag{3.2}$$

$$z = z_{b2}: \begin{cases} \partial_z L^a_\mu = 0, R^{1,2}_\mu = 0 \\ \partial_z(g_5 B_\mu + \tilde{g}_5 R^3_\mu) = 0, \tilde{g}_5 B_\mu - g_5 R^3_\mu = 0, \\ L^a_5 = 0, R^a_5 = 0, B_5 = 0 \end{cases} \tag{3.3}$$

where $g_5$ and $\tilde{g}_5$ are the five-dimensional gauge coupling for $SU(2)_{L,R}$ and $U(1)_{B-L}$, respectively. The motivation for these boundary conditions is given in Ref. [9]. These boundary conditions are consistent with models that have a bifundamental Higgs field in the $(2, 2)_0$ representation located at $z = 0$ that break $SU(2)_L \times SU(2)_R \rightarrow SU(2)_D$. On the second IR brane, at $z = z_{b2}$, there is a boundary

[2]We will now use $r_1$ instead of $r$ to denote the distance of the first brane to the origin. Also, we will only consider half of the space for most of the discussion since the other half is obtained by orbifolding about the origin.

Higgs in the $(1,2)_{1/2}$ representation under $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, which will break $SU(2)_R \times U(1)_{B-L} \to U(1)_Y$. In order to decouple the boundary Higgs fields from the theory, we take all of the Higgs vacuum expectation values to infinity. In addition to the boundary conditions, we impose continuity for the wave function and their derivatives at the UV brane, located at $z = z_b$, as dictated by the equations of motion. The bulk equation of motion for the gauge fields is

$$\left[\partial_{z'}^2 - \frac{1}{z'}\partial_{z'} + \frac{q^2}{k_{1,2}^2}\right]\psi(z') = 0, \tag{3.4}$$

where $z' = -k_1 z + 1$ or $k_2 z + C$ for $0 \le z \le z_b$ and $z_b \le z \le z_{b2}$, respectively. The solution to this equation is given by

$$\psi_i^d = \begin{cases} (-k_1 z + 1)(a_i^d J_1(q_i(-z + 1/k_1)) + b_i^d Y_1(q_i(-z + 1/k_1))), & 0 \le z \le z_b \\ (k_2 z + C)(a_i'^d J_1(q_i(z + C/k_2)) + b_i'^d Y_1(q_i(z + C/k_2))), & z_b \le z \le z_{b2} \end{cases}, \tag{3.5}$$

where $d$ labels the corresponding gauge bosons ($W_\pm, L3, B, R3$). Following [9], we expand the fields in their Kaluza-Klein modes as follows:

$$B_\mu(x, z) = \frac{1}{\tilde{g}_5} a_0 \gamma(x) + \sum_{j=1}^{\infty} \psi_j^B(z) Z_\mu^j(x), \tag{3.6}$$

$$L_\mu^3(x, z) = \frac{1}{g_5} a_0 \gamma(x) + \sum_{j=1}^{\infty} \psi_j^{L3}(z) Z_\mu^j(x), \tag{3.7}$$

$$R_\mu^3(x, z) = \frac{1}{g_5} a_0 \gamma(x) + \sum_{j=1}^{\infty} \psi_j^{R3}(z) Z_\mu^j(x), \tag{3.8}$$

$$L_\mu^\pm(x, z) = \sum_{j=1}^{\infty} \psi_j^{L_\pm}(z) W_\mu^{j\pm}(x), \tag{3.9}$$

$$R_\mu^\pm(x, z) = \sum_{j=1}^{\infty} \psi_j^{R_\pm}(z) W_\mu^{j\pm}(x). \tag{3.10}$$

### A. Oblique corrections

In order to calculate the electroweak corrections in our model we ensure that all corrections are oblique. This is done by adjusting the coupling of the fermions localized at $z = z_{b2}$ so that the zero mode couplings are equal to the SM couplings at tree level. For our model the relations are

$$-\frac{\tilde{g}_5 \psi_1^{(B)}(z_{b2})}{g_5 \psi_1^{(L3)}(z_{b2})} = \frac{g'^2}{g^2}, \tag{3.11}$$

$$g_5 \psi_1^{(L\pm)}(z_{b2}) = g, \tag{3.12}$$

$$g_5 \psi_1^{(L3)}(z_{b2}) = g \cos\theta_W. \tag{3.13}$$

For the photon kinetic term, we canonically normalize it as follows:

$$Z_\gamma = ((a_0/\tilde{g}_5)^2 + (a_0/g_5)^2) I = 1, \tag{3.14}$$

$$I = \int_{-z_{b2}}^{z_{b2}} e^{-A(z)/2} dz. \tag{3.15}$$

Equations (3.12) and (3.13) are used to determine the correct normalization for the $W$ and $Z$ wave functions.

Given the gauge field's wave functions, we calculated the oblique corrections using the relations between the vacuum polarization and the wave function renormalization $Z_\gamma = 1 - \Pi'_{QQ}$, $Z_W = 1 - g^2\Pi'_{11}$, and $Z_Z = 1 - (g^2 + g'^2)\Pi'_{33}$ [18]. The wave function renormalizations are give by

$$\begin{aligned} Z_W &= \int_{-z_{b2}}^{z_{b2}} [\psi^W]^2 e^{-A(z)/2} dz \\ &= \int_{-z_{b2}}^{z_{b2}} ([\psi^{L_+}]^2 + [\psi^{R_+}]^2) e^{-A(z)/2} dz, \end{aligned} \tag{3.16}$$

$$\begin{aligned} Z_Z &= \int_{-z_{b2}}^{z_{b2}} [\psi^Z]^2 e^{-A(z)/2} dz \\ &= \int_{-z_{b2}}^{z_{b2}} ([\psi^{L3}]^2 + [\psi^{R3}]^2 + [\psi^B]^2) e^{-A(z)/2} dz, \end{aligned} \tag{3.17}$$

and the zero momentum vacuum polarizations are

$$\Pi_{11}(0) = \frac{1}{g^2} \int_{-z_{b2}}^{z_{b2}} ([\partial_z \psi^{L_+}]^2 + [\partial_z \psi^{R_+}]^2) e^{-A(z)/2} dz, \tag{3.18}$$

$$\begin{aligned} \Pi_{33}(0) = \frac{1}{g^2 + g'^2} \int_{-z_{b2}}^{z_{b2}} &([\partial_z \psi^{L3}]^2 + [\partial_z \psi^{R3}]^2 \\ &+ [\partial_z \psi^B]^2) e^{-A(z)/2} dz. \end{aligned} \tag{3.19}$$

The Peskin-Takeuchi oblique corrections as a function of vacuum polarization are defined as [18]

$$S = 16\pi(\Pi'_{33} - \Pi'_{3Q}), \tag{3.20}$$

$$T = \frac{4\pi}{\sin^2\theta_W \cos^2\theta_W M_Z^2}(\Pi_{11}(0) - \Pi_{33}(0)), \tag{3.21}$$

$$U = 16\pi(\Pi'_{11} - \Pi'_{33}). \tag{3.22}$$

Since we are only considering the tree level corrections, $\Pi'_{3Q} = 0$. As an input to our model, we use the values of the SM electroweak parameters at the $Z$ pole: $M_W = 80.045$ GeV, $\sin^2\theta_W = 0.231$, and $\alpha = 127.9$. We also assume $k_1 r_1 = 30$. In the limit $r_2 \to 0$, $M_W$ sets the size of the extra dimension to be $r_1 = 68.5\ \text{TeV}^{-1}$. Since this is the limit of the standard Higgsless model, we find $T = U = 0$ and $S \sim 6\pi/(g^2(k_1 r_1)) \sim 1.4$ as in [9]. Since we are only interested in showing that the $S$ parameter decreases while preserving $T \sim 0$ and unitarity, we do not do a complete survey of the parameter space. For our analysis we set $k_2$ to be equal to the value of $k_1$ in the $r_2 \to 0$ limit. As we increase $r_2$, we find $r_1$ decreases in order to produce the proper $M_W$. Figure 4 shows the behavior of the $S$ parameter as we increase $r_2$. We find the $S$ parameter decreases. For $r_2 = 60$ we also checked that the lightest $W$ and $Z$ excitations are less than 1800 GeV and therefore unitarity is preserved [9]. This provides another mechanism for lowering the $S$ parameter in addition to including brane kinetic terms [9] and bulk fermions [19].

### B. The dual CFT description

A dual description using the AdS/conformal field theory (CFT) correspondence would include two CFTs coupled in the UV. There are two strongly interacting sectors in the IR: one is responsible for breaking $SU(2)_R \times U(1)_{B-L} \to U(1)_Y$, and one is responsible for breaking $SU(2)_L \times SU(2)_R \to SU(2)_D$. In contrast to the original Higgsless model in which the gauge group breaks to the standard model in the UV, our model breaks to the standard model at an IR scale that changes as we vary $r_2$. We find that as this scale decreases, the $S$ parameter also decreases.

The two strongly interacting scales become comparable when $r_2 \approx 65\ \text{TeV}^{-1}$, which corresponds to $k_1 r_1 \approx k_2 r_2$. At this scale, we find additional light resonances corresponding to the extended gauge group. The appearance of the resonances would affect four-Fermi operators and can be used to constrain our model. However, for the parameter space we consider in Fig. 4, the next lightest resonance above the $W$ and $Z$ mass is never below 1 TeV.

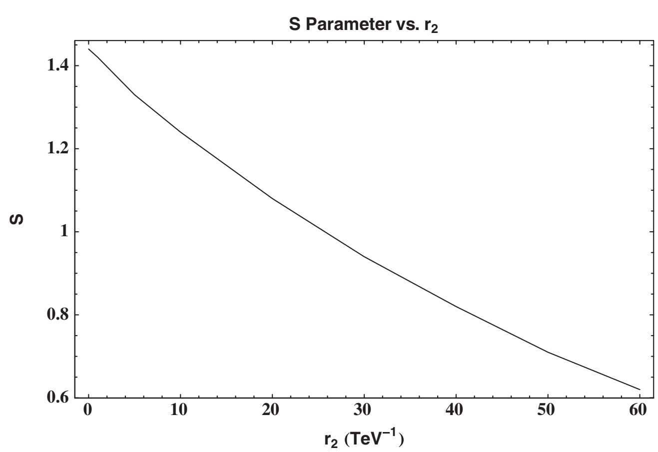


FIG. 4. Plot of the $S$ parameter as a function of $r_2$.

It would be interesting to understand from the four-dimensional perspective why the $S$ parameter decreases as a result of the additional strongly interacting sector. Future work could explore this interpretation as well as survey the parameter space keeping the constraint on four-Fermi operators in mind.

## IV. CONCLUSIONS

In the first section, we presented a model that is a hybrid between RSI and RSII. The model has a negative tension brane located at an orbifold fixed point ($y = 0$) and two identical positive tension branes located at $y = \pm r$. The fifth dimension extends to infinity as in RSII; however, the presence of the positive tension branes produces graviton resonances, which coincide with the discrete RSI spectrum. This model is attractive since it both solves the hierarchy problem and produces a continuum of KK graviton modes. As in both the RSI and RSII models, four-dimensional gravity can be recovered. Stability of our model is ensured by placing the negative tension brane at an orbifold fixed point.

In the second section of the paper, negative tension branes were brought in from infinity to cut off the space at an orbifold fixed point. We included $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ fields in the bulk and broke to the standard model on the far brane. The distances between the branes are scaled as to produce the correct $W$ mass. As in standard Higgsless electroweak symmetry breaking models, a large $S$ parameter along with vanishing $T$ and $U$ parameters were found when the second slice of our space was shrunk to zero. As the second slice of our space was increased, the $S$ parameter was lowered, while corrections to both $T$ and $U$ remained suppressed. We also find the lightest $W$ and $Z$ excitations stayed below 1800 GeV and therefore preserve unitarity. In conjunction with using brane kinetic terms and placing fermions in the bulk, this could be used as a useful mechanism for lowering the $S$ parameter.

Future work on these models could include trying to incorporate both Higgsless electroweak symmetry breaking and solutions to the hierarchy problem into a single model. It would also be interesting to explore how this model compares to other known mechanisms used to lower the $S$ parameter. In addition, an understanding of why the $S$ parameter decreases in the dual CFT description would be interesting.

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