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Brane-localized Kinetic Terms in the Randall-Sundrum Model [*] [†]

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Abstract

We examine the effects of boundary kinetic terms in the Randall-Sundrum model with gauge fields in the bulk. We derive the resulting gauge Kaluza-Klein (KK) state wavefunctions and their corresponding masses, as well as the KK gauge field couplings to boundary fermions, and find that they are modified in the presence of the boundary terms. In particular, for natural choices of the parameters, these fermionic couplings can be substantially suppressed compared to those in the conventional Randall-Sundrum scenario. This results in a significant relaxation of the bound on the lightest gauge KK mass obtained from precision electroweak data; we demonstrate that this bound can be as low as $m_1 \gtrsim 5$ TeV. Due to the relationship between the lightest gauge KK state and the electroweak scale in this model, this weakened constraint allows for the electroweak scale to be near a TeV in this minimal extension of the Randall-Sundrum model with bulk gauge fields, as opposed to the conventional scenario.

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1 Introduction

The Randall-Sundrum (RS) model [1] offers a new approach to the hierarchy problem. Within this scenario, the disparity between the electroweak and Planck scales is generated by the curvature of a 5-dimensional (5-d) background geometry which is a slice of anti-de Sitter (AdS_5) spacetime. This slice is bounded by two 3-branes of equal and opposite tensions sitting at the fixed points of an S^1/Z_2 orbifold which are located at $\phi = 0, \pi$ where ϕ is the coordinate of the 5^{th} dimension. The 5-d warped geometry then induces an effective 4-d scale Λ_π of order a TeV on the brane at $\phi = \pi$, the so-called TeV brane. Λ_π is exponentially smaller than the effective scale of gravity given by the reduced Planck scale, $\overline{M}_{Pl}$, with the suppression being determined by the product of the 5-d curvature parameter k and the separation r_c of the two branes, *i.e.*, $\Lambda_\pi = \overline{M}_{Pl} e^{-kr_c\pi}$. In this theory, the original parameters of the 5-d action are all naturally of the size $\sim \overline{M}_{Pl}$, while in the 4-d picture a hierarchy with $\Lambda_\pi/\overline{M}_{Pl} \lesssim 10^{-15}$ appears. It has been demonstrated [2] that this scenario is naturally stabilized without the introduction of fine-tuning for $kr_c \simeq 11$, which is the numerical value required to generate the hierarchy. This model leads to a distinct set of phenomenological signatures that may soon be revealed by experiments at the TeV scale; these have been examined in much detail [3, 4, 5].

For model building purposes, it is advantageous to extend the original framework of the RS model, where gravity alone propagates in the extra dimension, to the case where at least some of the Standard Model (SM) fields are present in the bulk [6]. The simplest such possibility is to have the SM gauge fields in the bulk [4, 5], while the other SM fields remain localized on the TeV brane. This scenario was examined in detail some time ago [4] where it was found that the couplings of the Kaluza-Klein (KK) excitations of the bulk gauge fields to matter on the boundary were enhanced by a factor of $\sqrt{2\pi kr_c}$ compared to those of the corresponding zero-mode states. Since $kr_c \simeq 11$ as discussed above, this enhancement is numerically significant and is approximately a factor of $\simeq 8.4$. Severe constraints are then imposed on the gauge KK excitations from precision electroweak measurements; these imply that the mass of the lightest gauge KK state must be heavy with $m_1 \geq 25 - 30$ TeV. Due to the relation in this model between the masses of the KK states and the scale of physics on the TeV brane, this bound then correspondingly requires that $\Lambda_\pi \geq 100$ TeV. This places such a scenario in a very unfavorable light in terms of its original motivation of resolving the hierarchy.

Here, we investigate whether Brane Localized Kinetic Terms (BLKT) for bulk gauge fields can modify these results. Recently, Carena, Tait and Wagner [7] examined the phenomenological influence of such localized kinetic terms for bulk gauge fields within the context of flat TeV^{-1}-sized extra dimensions. These authors showed that such terms can lead to significant modifications to the gauge KK spectrum, as well as to the KK state self-couplings and couplings to fields remaining on orbifold fixed points. Localized kinetic terms are expected to be present on rather general grounds in any orbifolded theory, as was shown by Georgi, Grant and Hailu [8], since they are needed to provide counterterms for divergences that are generated at one loop. In addition, these boundary terms can have important implications in a number of different situations such as model building [9], the construction of GUTs in higher dimensions with flat [10] or warped [11] geometries, and, in a different context, the quasilocalization of gravity [12]. Up to now the phenomenological implications of BLKT in the case of warped extra dimensions have not been examined.

We examine the effects of localized kinetic terms in the case where the SM gauge fields are present in the RS bulk and the remaining SM fields are on the TeV-brane. As in the case of TeV^{-1}-sized extra dimensions, we find that these boundary terms alter both the KK spectrum as well as their associated couplings to the remaining wall fields, but in a manner which is quite distinctive due to the curved geometry in this scenario. We will show that for a reasonable range of model parameters these modifications can be quite significant and lead to a natural reduction in the strength of the KK couplings. This allows the restrictions on Λ_π to be reduced by almost an order of magnitude and it would then no longer be unfavorable to have gauge fields be the only SM fields present in the RS bulk. The next section contains a detailed discussion of the modification of the usual gauge boson analysis within the RS framework due to the existence of brane localized kinetic terms. In Section 3 we discuss the implication of these results for the phenomenology of this scenario and our conclusions can be found in Section 4. Special limits of the KK mass equation are discussed in an Appendix.

2 Formalism

In this section, we derive the wavefunctions and the KK mass eigenvalue equation for bulk gauge fields with BLKT in the RS scenario. We then compute the coupling of the gauge field KK modes to 4-d fermions localized on the TeV brane, where the scale of physics is of order the electroweak scale. Our derivation will demonstrate how the boundary terms modify these couplings from their original form and result in a possible softening of the precision electroweak bounds on the mass of the lightest KK gauge state. In the rest of the paper, KK modes refer to those from the bulk gauge fields, unless otherwise specified.

We perform our calculations for a $U(1)$ gauge field; the generalization to the case of non-Abelian fields is straightforward. Let us start with the 5-d action for the gauge sector

$$\begin{aligned} S_A = \; & - \; \frac{1}{4}\int d^5x\,\sqrt{-G}\,\Big\{G^{AM}G^{BN}F_{AB}F_{MN} \\ & + \; [c_0\,\delta(\phi) + c_\pi\,\delta(\phi-\pi)]\,G^{\alpha\mu}G^{\beta\nu}F_{\alpha\beta}F_{\mu\nu}\Big\}\,, \end{aligned} \tag{1}$$

where the terms proportional to c_0 and c_π represent the kinetic terms for the branes located at $\phi = 0$, π, respectively, and the metric G_{MN} is given by the RS line element

$$ds^2 = e^{-2\sigma}\eta_{\mu\nu}dx^\mu dx^\nu - r_c^2 d\phi^2 \quad ; \quad \sigma = kr_c|\phi| \,. \tag{2}$$

Here the dimensionless constants $c_{0,\pi}$ would naturally be expected to be of order unity. Note that we can neglect terms proportional to $\partial_5 A_\mu$ on either brane as demonstrated by [7]. We will work in the $A_5 = 0$ gauge so that terms including A_5 can also be neglected. Thus only the usual 4-d kinetic term appears in the brane contributions to the action. In our notation, $A, B, \ldots = 0, \ldots, 3, 5$ and $\alpha, \beta, \ldots = 0, \ldots, 3$. Here, k is the 5-d curvature scale, r_c is the compactification radius, and $x^5 = r_c\phi$ with $-\pi \le \phi \le \pi$. The field strength is given by

$$F_{MN} = \partial_M A_N - \partial_N A_M \,. \tag{3}$$

The gauge field is assumed to have the KK expansion

$$A_\mu(x, \phi) = \sum_n A_\mu^{(n)}(x) \frac{\chi^{(n)}(\phi)}{\sqrt{r_c}} \,. \tag{4}$$

Substituting the above expansion in the action (1) yields

$$\begin{aligned} S_A = & - \frac{1}{4r_c}\int d^5x \sum_{m,n} \Big\{ [1 + c_0\,\delta(\phi) + c_\pi\,\delta(\phi - \pi)]\, F^{\mu\nu(m)} F_{\mu\nu}^{(n)} \chi^{(m)}\chi^{(n)} \\ & - (2/r_c^2)\, A^{\mu(m)} A_\mu^{(n)} e^{-2\sigma} [(d/d\phi)\chi^{(m)}][(d/d\phi)\chi^{(n)}] \Big\} \,, \end{aligned} \tag{5}$$

where now all the Lorentz indices are contracted using the 4-d Minkowski metric. In order to cast this action in the diagonal form, we demand

$$\int d\phi\, [1 + c_0\,\delta(\phi) + c_\pi\,\delta(\phi - \pi)]\, \chi^{(m)}\chi^{(n)} = Z_n \delta^{mn} \tag{6}$$

and

$$\frac{1}{r_c^2}\int d\phi\, e^{-2\sigma} \Big[\frac{d}{d\phi}\chi^{(m)}\Big]\Big[\frac{d}{d\phi}\chi^{(n)}\Big] = Z_n m_n^2\, \delta^{mn} \,. \tag{7}$$

These equations imply

$$\frac{d}{d\phi}\left(e^{-2\sigma}\frac{d}{d\phi}\chi^{(n)}\right) + [1 + c_0\,\delta(\phi) + c_\pi\,\delta(\phi-\pi)]\,r_c^2\,m_n^2\,\chi^{(n)} = 0\,. \tag{8}$$

Away from the boundaries at $\phi = 0, \pi$, the solution to Eq.(8) is given by [4, 5]

$$\chi^{(n)}(\phi) = \frac{e^{\sigma}}{N_n}\left[J_1(z_n) + \alpha_n\,Y_1(z_n)\right], \tag{9}$$

where $z_n(\phi) \equiv (m_n/k)e^{\sigma}$; J_q and Y_q denote the usual Bessel functions of order q. The normalization N_n is set by the orthonormality condition

$$\int_{-\pi}^{+\pi} d\phi\;\chi^{(m)}\chi^{(n)} = \delta^{mn}\,, \tag{10}$$

implying that the zero mode is given by $\chi^{(0)} = 1/\sqrt{2\pi}$. Integrating Eq.(8) around the fixed points $\phi = 0$ and $\phi = \pm\pi$ yields

$$2\frac{d}{d\phi}\chi^{(n)}(0) + c_0\,r_c^2 m_n^2\,\chi^{(n)}(0) = 0 \tag{11}$$

and

$$2\frac{d}{d\phi}\chi^{(n)}(\pi) + e^{2kr_c\pi}c_\pi r_c^2\,m_n^2\,\chi^{(n)}(\pi) = 0\,, \tag{12}$$

respectively.

Let $\varepsilon_n \equiv z_n(0)$. For light KK modes, with masses $m_n \sim 1$ TeV, which are of phenomenological interest, $\varepsilon_n \ll 1$; for $k \sim \overline{M}_{Pl}$ we have $\varepsilon_n \sim 10^{-15}$. Substituting for $\chi^{(n)}$ from Eq.(9) in Eq.(11), we obtain

$$\alpha_n = -\frac{(1+\delta_0\varepsilon_n^2)J_1(\varepsilon_n) + \varepsilon_n J_1'(\varepsilon_n)}{(1+\delta_0\varepsilon_n^2)Y_1(\varepsilon_n) + \varepsilon_n Y_1'(\varepsilon_n)}\,, \tag{13}$$

or, using the standard Bessel function relations described in the Appendix

$$\alpha_n = -\frac{J_0(\varepsilon_n) + \delta_0 \varepsilon_n J_1(\varepsilon_n)}{Y_0(\varepsilon_n) + \delta_0 \varepsilon_n Y_1(\varepsilon_n)}\,, \tag{14}$$

where $\delta_0 \equiv c_0 k\, r_c/2$ and a prime ($\prime$) denotes a derivative. Note that with c_0 of order unity, we may expect values of δ_0 of order 10 since $kr_c \sim 11$ to solve the hierarchy. Similarly, from Eq.(12), we obtain

$$(1 + \delta_\pi x_n^2) J_1(x_n) + x_n J_1'(x_n) + \alpha_n[(1 + \delta_\pi x_n^2) Y_1(x_n) + x_n Y_1'(x_n)] = 0\,. \tag{15}$$

Again employing the usual Bessel function identities and dividing by x_n, this simplifies to

$$J_0(x_n) + \delta_\pi x_n J_1(x_n) + \alpha_n[Y_0(x_n) + \delta_\pi x_n Y_1(x_n)] = 0\,, \tag{16}$$

where $\delta_\pi \equiv c_\pi k\, r_c/2$ and $x_n \equiv z_n(\pi)$. The roots x_n of Eq.(16) yield the KK mass spectrum, $m_n = x_n k e^{-kr_c\pi}$. We will study various limits of this root equation in the Appendix. Following the same arguments as above we may expect, assuming no fine-tuning, $|\delta_{0,\pi}| \lesssim 10$.

For exponentially small values of ε_n, it may first appear that the δ_0-dependent corrections to α_n in Eq.(13) are suppressed by powers of ε_n. However, an expansion in the small parameter ε_n reveals that this is not the case. We find in this limit

$$\alpha_n \simeq -\frac{\pi/2}{\ln(x_n/2) - k\, r_c\, \pi + \gamma - \delta_0}\,, \tag{17}$$

where $\gamma \approx 0.577$ is Euler's constant. α_n remains small over most regions of interest here, except for when $\delta_0 \simeq -kr_c\pi \simeq -35.5$. There are two sources of modifications to α_n due to the presence of the boundary terms, as compared to the original RS framework: (1) the appearance of the term proportional to δ_0 in the denominator which is not suppressed by the exponential warp factor, and (2) the BLKT-modified values of x_n. In our numerical analyses

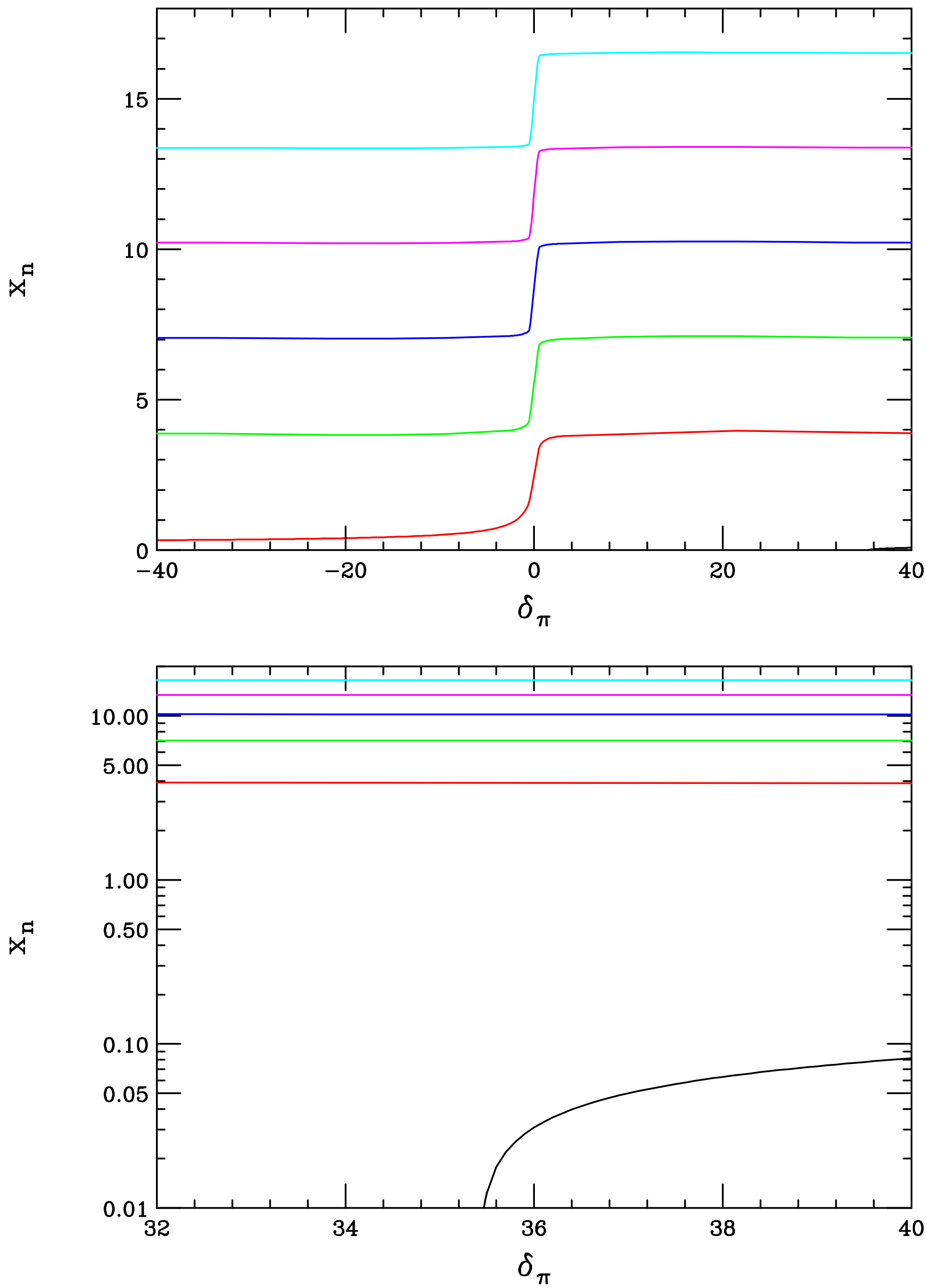

Figure 1: (a) Top panel. The behavior of the first five roots as δ_π is varied, assuming $\delta_0 = 0$. Note the sharp increase in x_n near $\delta_\pi = 0$ for all KK levels. (b) Lower panel. The large δ_π region is expanded to show the new root originating at the value $\delta_\pi = kr_c\pi \simeq 35.5$.

we will use the exact expression for α_n, rather than the approximate version given above, and we take $kr_c = 11.27$.

To get a flavor of how the KK tower mass spectrum is modified in this scenario, we must solve Eq.(16) for a range of values for $\delta_{0,\pi}$. We first take the simplest possibility by setting $\delta_0 = 0$, and examine the roots x_n as a function of δ_π. Our results are displayed in Fig. 1. From Fig. 1a, we see that in the region near $\delta_\pi = 0$ the roots undergo a substantial shift but remain relatively independent of variations in δ_π away from the origin. For large positive values of δ_π, the roots are almost the same as those obtained in the case of the graviton KK spectrum in the original RS framework [3]. This can be seen analyticallly by taking the limit of large δ_π in Eq.(16) and neglecting $\alpha_n \ll 1$. As we will see below, the region $\delta_\pi < -1/2$ is unphysical. When the value of δ_π exceeds $kr_c\pi$, a new root appears which is barely perceptible in Fig. 1a, but whose turn-on is shown in detail in Fig. 1b. The analytical source of this root is discussed in the Appendix. Physically, it is similar to the 'collective' mode, which appears in the limit where the brane coupling goes to infinity, in the case of two branes being present in a TeV^{-1}-sized extra dimension as discussed by Carena *et al.* [7]. Here, due to the exponential warp factor, this limit is reached for rather modest values of δ_π.

When $\delta_0 \neq 0$ the modifications to the roots are sensitive to the sign of δ_0. However, in the parameter range of interest to us here, the roots appear to be well described by the results shown in Fig. 1 at the level of a few %. The major difference in this case is the location of where the additional root materializes; as shown in the Appendix, this root now appears at the value $\delta_\pi = kr_c\pi + \delta_0$. When $\delta_0 > (<)0$ the turn-on of this root moves to larger (smaller) values of δ_π. In the case $\delta_0 = -\delta_\pi$, Fig. 2 shows that the new root appears at $\delta_\pi = kr_c\pi/2$ as expected. In order to avoid the phenomenologically troublesome regions of parameter space where the new root appears and results in a very light KK gauge state, we will restrict our analysis to the range $\delta_\pi < kr_c\pi + \delta_0$ in what follows.

When $\delta_\pi = 0$, the roots exhibit a weak dependence on δ_0 as shown in Fig. 3. Away from the $\delta_0 = -kr_c\pi$ region, the δ_0 dependence of the roots is essentially isolated in the coefficients α_n as seen in Eq.(17). From this we see that the roots have a slightly negative slope with increasing δ_0.

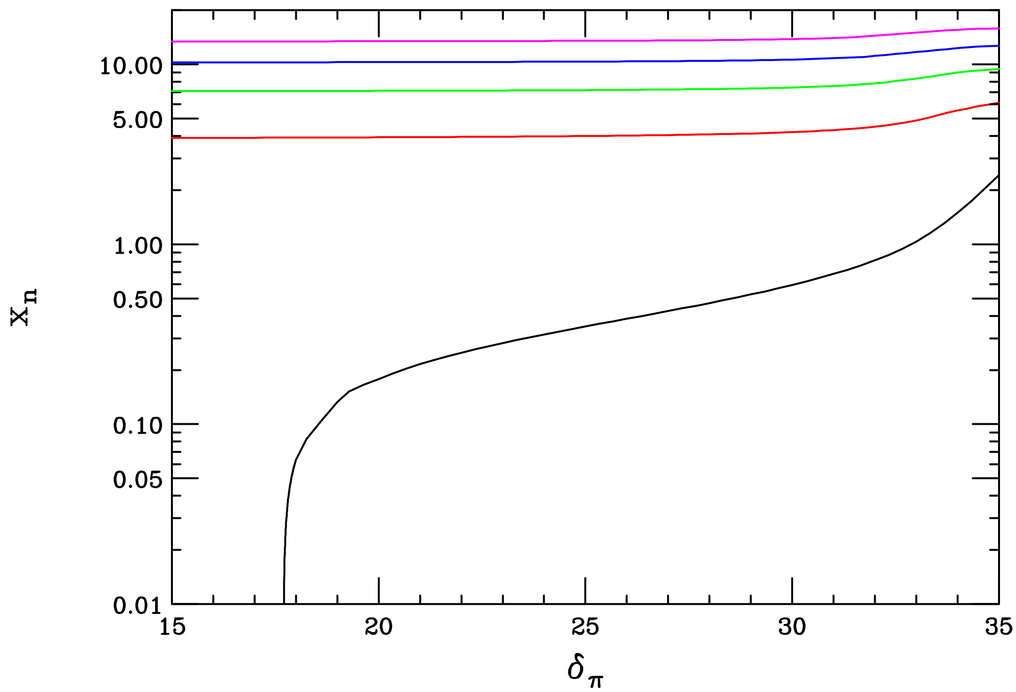

Figure 2: The large δ_π region for the case $\delta_0 = -\delta_\pi$. Here, the new root appears at the value $\delta_\pi = kr_c\pi/2 \simeq 17.7$.

3 Analysis

In this section, we study the couplings of the KK modes to the boundary fermions at $\phi = \pi$. The values of these couplings are the driving force behind the important constraints on the mass of the lightest KK gauge field in the original RS scenario [4]. We would like to reexamine these bounds in the presence of BLKT. To begin, we note that the diagonalization

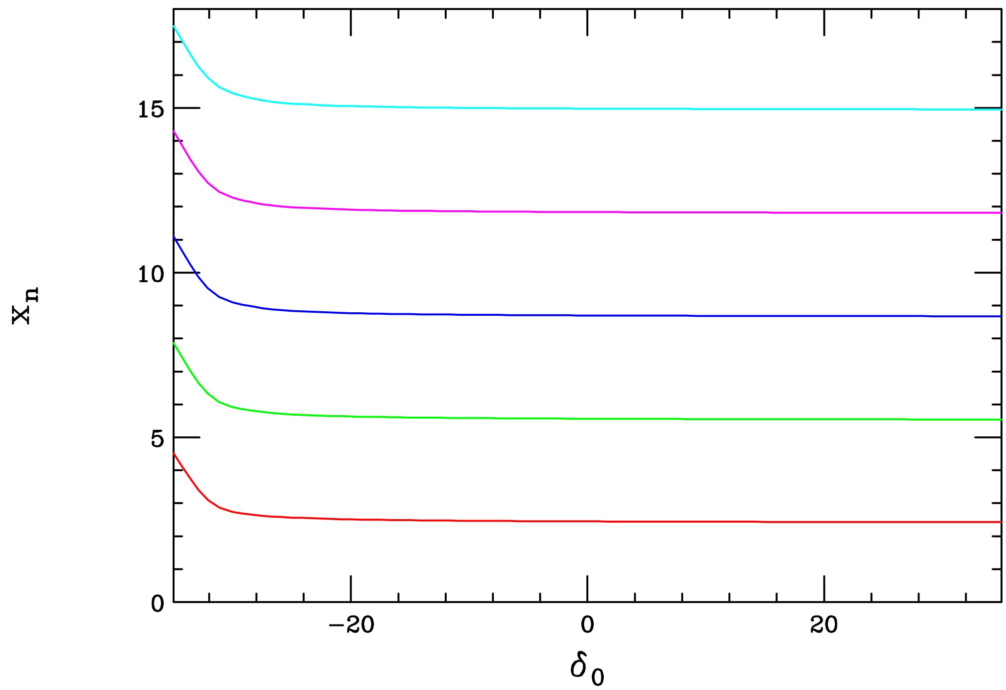

Figure 3: The behavior of the first five roots as a function of δ_0 taking $\delta_\pi = 0$.

conditions (6) and (7) yield the 4-d action

$$S_A = -\frac{1}{4}\int d^4x \sum_n Z_n \left(F^{\mu\nu(n)}F^{(n)}_{\mu\nu} - 2m_n^2 A^{\mu(n)}A^{(n)}_\mu\right) , \tag{18}$$

where we have

$$Z_n = 1 + c_0[\chi^{(n)}(0)]^2 + c_\pi[\chi^{(n)}(\pi)]^2 . \tag{19}$$

Note that for physical fields, we demand that $Z_n > 0$ for all $n \geq 0$.

The coupling of a bulk gauge field A_μ to fermions ψ localized at $\phi = \pi$ is given by the action

$$S_\psi = -\int d\phi\, d^4x\, (detV)\, g_5\, \bar{\psi} V^\mu_\alpha \gamma^\alpha A_\mu(x,\phi)\psi\, \delta(\phi - \pi) , \tag{20}$$

where $V^\mu_\alpha = e^\sigma \eta^\mu_\alpha$ is the vielbein, $detV = e^{-4\sigma}$, and g_5 is the 5-d coupling constant. Using the expansion (4) in Eq.(20), we obtain

$$S_\psi = -\int d^4x \left[\frac{g_5}{\sqrt{2\pi r_c}} \bar{\psi}\gamma^\mu A_\mu^{(0)}\psi + \sum_{n\neq 0} g_5\sqrt{k}\bar{\psi}\gamma^\mu A_\mu^{(n)}\psi \right] , \tag{21}$$

where we have performed the redefinition $\psi \to e^{3kr_c\pi/2}\psi$ to make the ψ kinetic terms canonical. Here, we have used

$$N_n \simeq \frac{e^{kr_c\pi}}{\sqrt{kr_c}} J_1(x_n) , \tag{22}$$

where we have ignored α_n which is of order $\sim 10^{-2}$ for the parameter values which are phenomenologically relevant. To bring the gauge field kinetic terms in Eq.(18) into the canonical form, we require

$$A_\mu^{(n)} \to \frac{A_\mu^{(n)}}{\sqrt{Z_n}} . \tag{23}$$

Eqs.(21) and (23) then yield

$$S_\psi = -\int d^4x \left[g_0\, \bar{\psi}\gamma^\mu A_\mu^{(0)}\psi + \sum_{n\neq 0} g_n\bar{\psi}\gamma^\mu A_\mu^{(n)}\psi \right] , \tag{24}$$

with $g_0 \equiv g_5/\sqrt{2\pi r_c Z_0}$ and $g_n \equiv g_5\sqrt{k/Z_n}$. Here, g_0 is identified as the usual 4-d coupling constant for the interactions of the fermions with the zero-mode KK state. The KK excitation couplings are given by g_n, which due to the presence of the boundary terms are now mode dependent. Without the presence of the brane kinetic terms, the original RS model yields [4, 5]

$$\left.\frac{g_n}{g_0}\right|_{RS} = \sqrt{2\pi k r_c} , \tag{25}$$

whereas, in the presence of the localized kinetic terms, we find that (as long as the value of

δ_0 is not close to $-kr_c\pi$ such that the terms proportional to $\alpha_n \simeq 10^{-2}$ may be neglected)

$$\left.\frac{g_n}{g_0}\right|_{BLKT} = \sqrt{2\pi kr_c}\sqrt{\frac{Z_0}{Z_n}}\;. \tag{26}$$

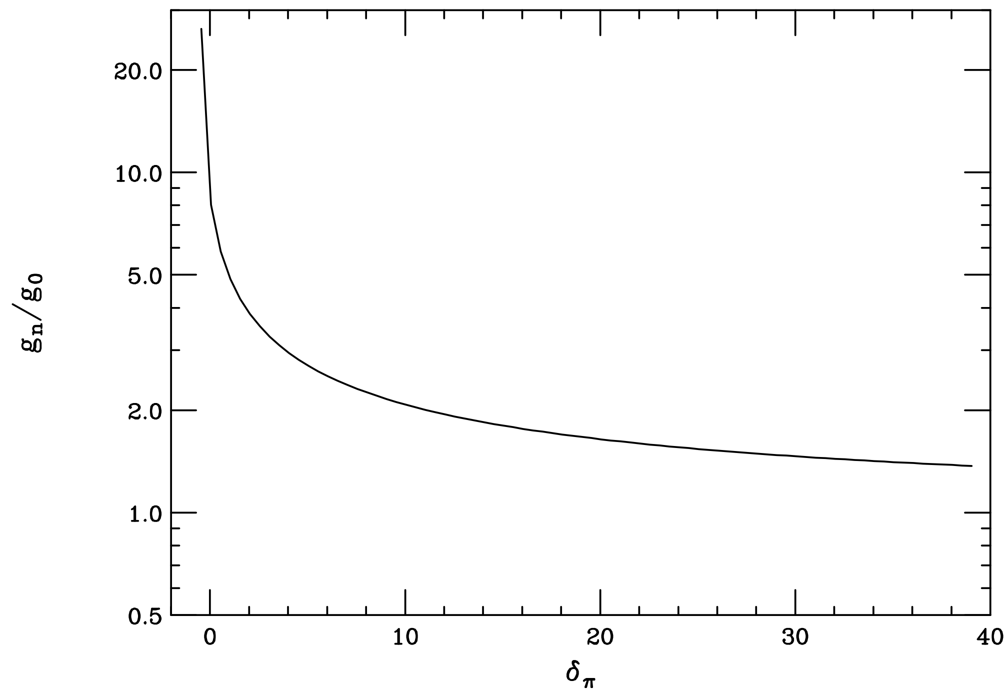

Figure 4: The ratio of the KK couplings to that of the zero-mode as a function of δ_π for the case $\delta_0 = 0$.

We can now make an estimate of the size of the effect of the brane localized kinetic terms on the ratio of the couplings in Eq.(26). Using Eq.(19), we obtain

$$Z_0 = 1 + \frac{c_0 + c_\pi}{2\pi}\,, \tag{27}$$

and for typical values of parameters

$$Z_n \simeq 1 + c_\pi kr_c + \frac{c_0 c_\pi^2 \alpha_n^2 (kr_c)^3}{\pi^2[J_0(x_n) + \alpha_n Y_0(x_n)]^2}\,. \tag{28}$$

Note that unlike in the case of TeV^{-1}-sized extra dimensions, here, to a very good approximation, Z_n is n-independent for small values of α_n. In the same approximation of small α_n, since we must have $Z_0, Z_n > 0$, these equations imply that $c_0 + c_\pi > -2\pi$, and similarly $c_\pi > -(kr_c)^{-1}$. Translating these to constraints on $\delta_{0,\pi}$ implies that $\delta_0 + \delta_\pi > -kr_c\pi \simeq -35.5$ and $\delta_\pi > -1/2$.

We now see how the n^{th}-mode KK gauge couplings compare to the usual RS scenario as $\delta_{\pi,0}$ are varied. For purposes of demonstration, we only consider the case $\delta_0 = 0$. Generalizations to non-zero values of δ_0 are straightforward. Fig. 3a displays the significant fall-off of the ratio g_n/g_0 as a function of δ_π. We see that substantial reduction over the original RS framework, that is the RS model without the presence of boundary terms, is possible.

Since $kr_c \sim 11$, for $c_0 \sim c_\pi \sim 1$, we obtain a typical value of $\sqrt{Z_0/Z_n} \approx 0.3$, resulting in a substantial suppression of the gauge KK mode coupling to the fermions localized on the TeV brane as compared to the original RS model. Hence, we expect that the inclusion of the localized kinetic terms can typically result in a significant loosening of the precision electroweak bounds [4] on the model. To make this more quantitative we need to consider the quantity

$$\tilde{V}(\delta_\pi, \delta_0) \equiv \sum_n \Big(\frac{g_n}{g_0}\Big)^2 \Big(\frac{x_1}{x_n}\Big)^2 , \tag{29}$$

where the x_n are the roots given above, and examine its detailed dependence on $\delta_{0,\pi}$. This quantity describes the set of contact interactions induced by dimension-6 operators arising from KK gauge exchanges in precision electroweak observables [4, 13]. Since for $\delta_{0,\pi} = 0$ we know that precision measurements [14] imply the bound $m_1|_{RS} \gtrsim 25$ TeV for the lightest gauge KK excitation [4], we can obtain the corresponding constraint for non-zero values of $\delta_{0,\pi}$. We obtain the bound $m_1|_{BLKT} \geq m_1|_{RS}[\tilde{V}(\delta_\pi, \delta_0)/\tilde{V}(0,0)]^{1/2}$; the numerical result of this analysis is presented in Fig. 5. Here, we see that the bound on m_1 falls rapidly with increasing δ_π. For $\delta_\pi \sim 10$, we see that m_1 can be as low as $5-6$ TeV which will lead to an observable signal in Drell-Yan production for the first gauge KK excitation at the LHC. This is demonstrated in Fig 6 for the case of $m_1 = 5$ TeV. In addition, since the bounds on m_1 are reduced here by a factor of $\simeq 5$ below those of the conventional RS framework, the resulting bound on the scale of physics on the TeV brane is also reduced to $\Lambda_\pi \gtrsim 20$. This removes a possible new hierarchy that one finds in the usual scenario [4].

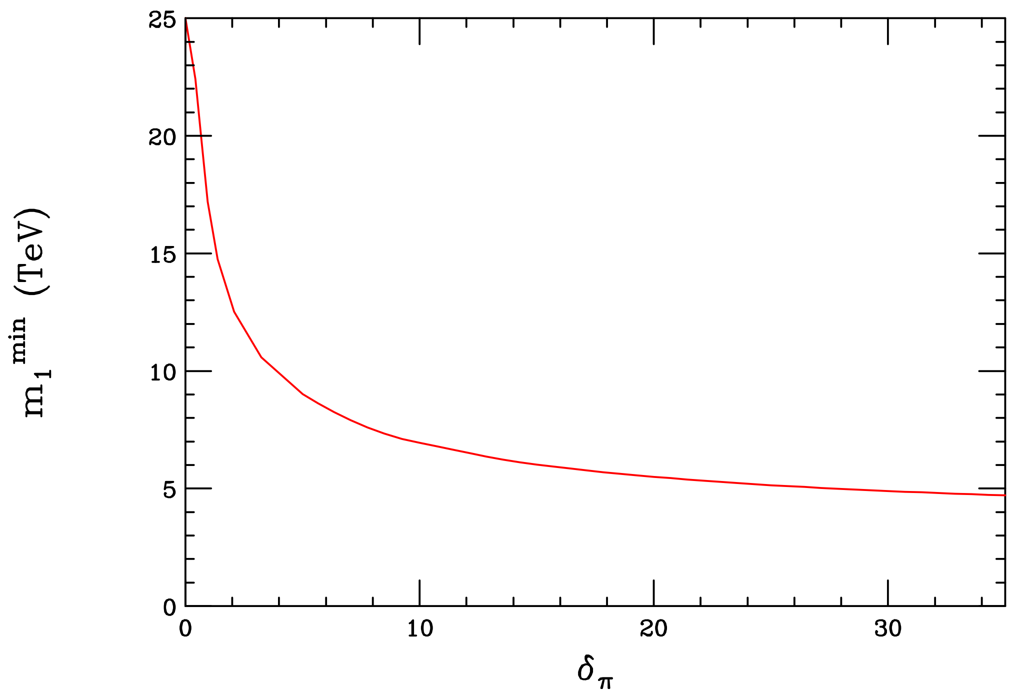

Figure 5: Lower bound on the mass of the lightest gauge boson KK excitation as a function of δ_π from a fit to the precision electroweak data for the case $\delta_0 = 0$.

4 Conclusions

A minimal, yet interesting, extension of the RS model is obtained by assuming that only gauge fields propagate in the bulk. However, this setup has been shown to be constrained by precision electroweak data to have a scale $\Lambda_\pi \gtrsim 100$ TeV for physics on the TeV brane [4]. This makes such a scenario less attractive as a means for resolving the gauge hierarchy. The stringent bound on the electroweak scale is a result of the strong coupling of the gauge KK

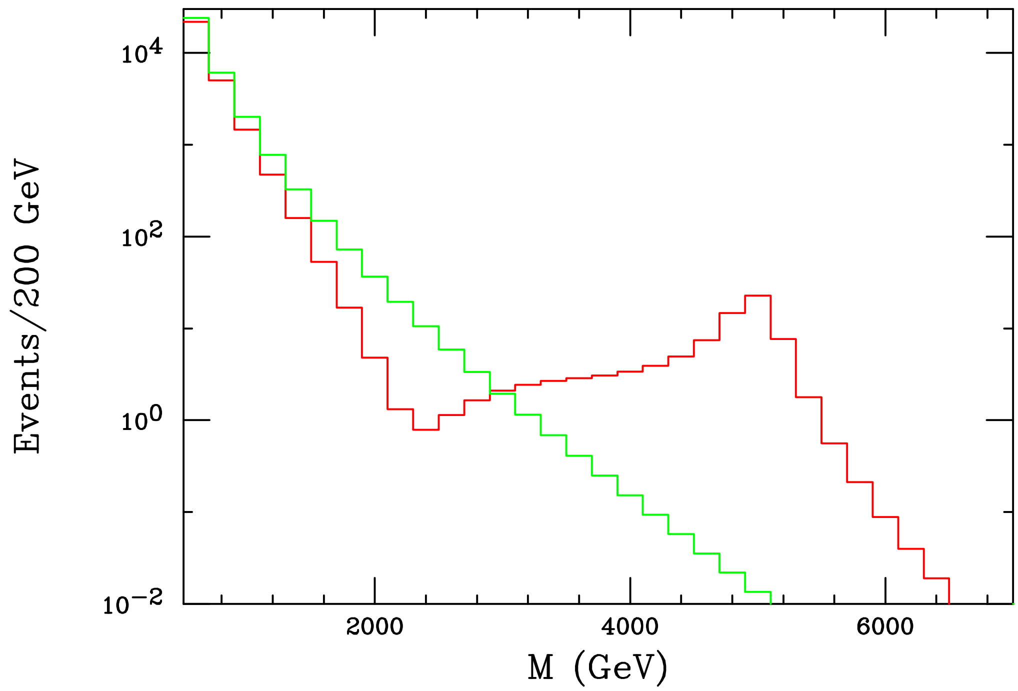

Figure 6: Drell-Yan event rate at the LHC for the first gauge KK state in the RS scenario with boundary terms. We have assumed a luminosity of 300 fb^{-1} and have set $m_1 = 5$ TeV. The steeply falling histogram represents the SM Drell-Yan continuum.

modes to the boundary fermions. It is thus interesting to examine if a modification of the boundary physics can lead to a relaxation of this constraint. This simple extension of the theory is also motivated by field theoretic considerations [8] is the addition of brane localized kinetic terms for the bulk gauge fields.

In this paper, we assumed the presence of gauge field localized kinetic terms on both the Planck and TeV branes present in the RS model. We then derived the wavefunctions and the mass spectrum of the gauge KK modes and the couplings of the KK states to the SM brane localized fermions were then calculated. It was shown that for natural choices of parameters, a substantial suppression of these KK couplings, compared to the case without BLKT, can be achieved. We then reexamined the precision electroweak bounds on the lightest gauge KK mass or, equivalently, Λ_π. We found that for a reasonable range of parameters, the mass of the lightest gauge KK field can be as low as $\simeq 5$ TeV and hence may be visible at the LHC. This weakening of the constraints from precision data results in a significantly less severe bound on Λ_π; one can now naturally obtain $\Lambda_\pi \sim 20$ TeV. This makes the RS framework more favorable, in the context of a solution to the hierarchy problem, thereby facilitating the construction of more realistic models based on the RS proposal.

Based on our observations in this paper, we expect that the addition of brane localized kinetic terms for other bulk fields, such as gravitons, could significantly alter the phenomenology of the RS model. This could also lead to the appearance of new features, such as the light KK mode that was discussed above, in the low energy theory.

Acknowledgements

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Appendix

Here, we will study the limiting behavior of Eq.(15). We will show that near special values of the boundary term coefficients δ_0 and δ_π, the equation develops a small root. We also discuss the $\delta_{0,\pi} \gg 1$ limit. For simplicity, we only consider varying the value of δ_π.

A Bessel function ζ obeys the following identity

$$x\,\zeta_0(x) = \zeta_1(x) + x\,\zeta_1'(x)\,, \tag{30}$$

where the subscripts denote the order of ζ. Using (30), we may rewrite Eq.(15) as

$$J_0(x_n) + \delta_\pi\, x_n J_1(x_n) + \alpha_n[Y_0(x_n) + \delta_\pi\, x_n Y_1(x_n)] = 0\,. \tag{31}$$

In the limit where $x_n \to 0$, Eqs.(31) and (17) yield

$$x_n^2 = \frac{4(\delta_\pi - \delta_0 - kr_c\pi)}{(kr_c\pi + \delta_0)(2\delta_\pi - 1) + (1 - \delta_\pi)} \,. \tag{32}$$

Eq.(32) is valid for $x_n \ll 1$ and we see that as $\delta_\pi \to (\delta_0 + kr_c\pi)^+$, $x_n^2 \to 0^+$, where the + superscript denotes an approach from above. Thus, for δ_π sufficiently close to the special value $\delta_0 + kr_c\pi$, a light mode appears in the spectrum. Note that in the original RS model with $\delta_0 = \delta_\pi = 0$, to get this small root, one has to require that $kr_c\pi \to 0^-$ in Eq.(32), which is not allowed.

We also note that for $\delta_0 > 1/2$, as $\delta_\pi \to \pm\infty$, Eq.(32) implies

$$x_n^2 \to \frac{2}{kr_c\pi + \delta_0 - 1/2} \,. \tag{33}$$

Since the validity of Eq.(32) requires $x_n \ll 1$, the limit (33) is a good approximation only for $|\delta_\pi| \gg \delta_0 \gg 1$. Hence, we see that in the limit $\delta_\pi \to \pm\infty$, $\delta_0 \gg 1$, a new light mode appears above the zero mode. In this limit, $\alpha_n \ll 1$, which implies that Eq.(31) is well-approximated by $J_1(x_n) = 0$, for $x_n \neq 0$. In this case, apart from the light mode given by (33), the rest of the gauge field KK modes have masses that are approximately degenerate with those of the KK gravitons from the original RS framework [3]. Our numerical studies support the conclusions reached in this Appendix. For example, Fig. 1b shows in the case $\delta_0 = 0$, that the new root begins at $\delta_\pi = kr_c\pi$ and is approaching the asymptotic value given by the above equation.

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