Constraints on the bulk standard model in the Randall-Sundrum scenario

Gustavo Burdman
Theoretical Physics Group, Lawrence Berkeley National Laboratory, Berkeley, California 94720
(Received 10 June 2002; published 16 October 2002)

We derive constraints on the Randall-Sundrum scenario with the standard model fields in the bulk. These result from tree level effects associated with the deformation of the zero mode wave functions of the W and the Z once electroweak symmetry is broken. Recently Csáki, Erlich and Terning pointed out that this implies large contributions to electroweak oblique parameters. Here we find that when fermions are allowed in the bulk the couplings of the W and the Z to zero-mode fermions are also affected. We perform a fit to electroweak observables, assuming universal bulk fermion masses and including all effects, and find constraints that are considerably stronger than for the case with fermions localized in the low energy boundary. These put the lowest Kaluza-Klein excitation out of reach of the CERN Large Hadron Collider. We then relax the universality assumption and study the effects of flavor violation in the bulk and its possible signatures.

DOI: 10.1103/PhysRevD.66.076003 PACS number(s): 11.10.Kk

I. INTRODUCTION

Theories with large extra dimensions have been recently introduced as an alternative framework to solve the hierarchy problem [1]. It is assumed that the geometry is factorizable and results in a product of Minkowski space with n compact dimensions. In this scenario, gravity propagates in the extra dimensions so that the strength of its coupling to matter confined in our four-dimensional world is determined by the scale $M_{\rm P}^2 = M^{n+2}V_n$, with M the fundamental scale of gravity and V_n the volume of the extra dimensions. In this way, the hierarchy between $M_{\rm P}$ and the weak scale results from the volume suppression, and the truly fundamental scale M can be of the order of 1 TeV.

An alternative scenario by Randall and Sundrum (RS) involves the use of a non-factorizable geometry in five dimensions [2]. The metric depends on the five-dimensional coordinate y and is given by

$$ds^2 = e^{-2\sigma(y)}\eta_{\mu\nu}dx^\mu dx^\nu - dy^2, \tag{1}$$

where x^μ are the four-dimensional coordinates, $\sigma(y) = k|y|$, with $k \sim M_P$ characterizing the curvature scale. The extra dimension is compactified on an orbifold S_1/Z_2 of radius r so that the bulk is a slice of AdS_5 space between two four-dimensional boundaries. The metric on these boundaries generates two effective scales: M_P and $M_P e^{-k\pi r}$. In this way, values of r not much larger than the Planck length can be used in order to generate a scale $\Lambda_r \simeq M_P e^{-k\pi r} \simeq 0(\mathrm{TeV})$, i.e. $kr \simeq (11-12)$, for generating the TeV on one of the boundaries.

In the original RS scenario, only gravity was allowed to propagate in the bulk, with the standard model (SM) fields confined to one of the boundaries. The inclusion of matter and gauge fields in the bulk has been extensively treated in the literature [3–9]. In this paper we are interested in examining the situation when the SM fields are allowed to propagate in the bulk. The exception in this so-called bulk SM is the Higgs field which must be localized on the TeV boundary in order for the W and the Z gauge bosons to get their observed masses [4]. As it was first noted in Ref. [7] the wave functions of the W and Z acquire a dependence on the fifth-dimensional coordinate due to the Higgs vacuum expectation value (VEV). Recently Csáki, Erlich and Terning [10] have studied the effects that result from this deformation of the zero modes in a scenario with only gauge fields in the bulk. They found large contributions to the oblique parameters S and T and the bound $\Lambda_r > 11$ TeV, which is slightly tighter than the ones previously obtained [8]. In this paper we consider a scenario where all or at least part of the fermions can propagate in the bulk. It is generally believed that this relaxes the bounds on Λ_r since the couplings of Kaluza-Klein excitations of gauge bosons to zero-mode fermions are not as strong as when fermion are confined to the TeV boundary. Here we show that there are additional effects resulting from the modified couplings of W and Z to the SM fermions that propagate in the bulk. Even when we consider these to be flavor universal, they result in non-oblique contributions to electroweak observables and in *stronger* constraints on Λ_r than the ones obtained with confined fermions. If flavor breaking in the bulk is allowed, then there are additional effects in flavor changing neutral processes.

In Sec. II we review the bulk SM and the existing bounds on the induced low energy scale Λ_r. In Sec. III we obtain the bounds on Λ_r coming from the deviation in the tree level couplings of fermions to W and Z, adding this effect to the contributions discussed in Ref. [10]. We derive these new constraints by making the simplifying assumption that the effects on the fermion couplings are flavor universal. In Sec. IV we study the effects of flavor violation in the bulk. Finally, in Sec. V we conclude.

II. THE BULK STANDARD MODEL

The five-dimensional action for bulk gauge fields is given by [3,4]

$$S_A = -\frac{1}{4}\int d^4x dy \sqrt{-g} F^{MN}F_{MN}, \tag{2}$$

where $g = \det(g_{MN}) = e^{-4\sigma(y)}$ and capital latin letters denote five-dimensional indexes. The field strength is in general written as

$$F_{MN} \equiv \partial_M A_N - \partial_N A_M + ig_5[A_M, A_N]. \tag{3}$$

We make the choice of gauge $A_y=0$. The Kaluza-Klein (KK) decomposition is given by

$$A_\mu(x,y)=\frac{1}{\sqrt{2\pi r}}\sum_{n=0}^{\infty} A_\mu^{(n)}(x)\chi^{(n)}(y). \tag{4}$$

Thus, the wave function of the gauge boson in the fifth dimension $\chi^{(n)}(y)$ obeys the differential equation

$$-\partial_y(e^{-2\sigma}\partial_y\chi^{(n)})=m_n^2\chi^{(n)}. \tag{5}$$

The solutions satisfy the normalization condition

$$\frac{1}{2\pi r}\int_{-\pi r}^{\pi r}dy \quad \chi^{(m)}\chi^{(n)}=\delta_{mn}, \tag{6}$$

and are

$$\chi^{(n)}(y)=\frac{e^\sigma}{N_n}\left[J_1\left(\frac{m_n}{\kappa}e^\sigma\right)+\alpha_n Y_1\left(\frac{m_n}{\kappa}e^\sigma\right)\right]. \tag{7}$$

In Eq. (7), N_n is the normalization constant derived from Eq. (6), and

$$\alpha_n=-\frac{J_0(m_n/k)}{Y_0(m_n/k)}, \tag{8}$$

where we defined

$$x_n\equiv\frac{m_n}{k} \quad e^{kr\pi}=\frac{m_n}{\Lambda_r}, \tag{9}$$

i.e. the mass of the KK excitation in units of the generated low energy scale. The zero mode is flat in the extra dimension: $\chi^{(0)}(y)=1$. Imposing continuity at $y=(0,\pi r)$ results in the condition [3,4]

$$J_0(x_n) \quad Y_0(x_n e^{-kr\pi})=J_0(x_n e^{-kr\pi}) \quad Y_0(x_n), \tag{10}$$

which determines the KK masses. For instance, for $e^{-kr\pi}=10^{-16}$, we have $x_1\simeq 2.45$, $x_2\simeq 5.6$, $x_3\simeq 8.70$, $x_4\simeq 11.8$, $x_5\simeq 15.0$, etc. with the KK masses given by Eq. (9).

The action for fermion fields in the bulk is given by [4,5]

$$S_f=\int d^4x dy \sqrt{-g}\left\{\frac{i}{2}\bar{\Psi}\hat{\gamma}^M[\mathcal{D}_M-\overleftarrow{D}_M]\Psi -\mathrm{sgn}(y)M_f\bar{\Psi}\Psi\right\}, \tag{11}$$

where the covariant derivative in curved space is

$$\mathcal{D}_M\equiv\partial_M+\frac{1}{8}[\gamma^\alpha, \quad \gamma^\beta]V_\alpha^N V_{\beta N;M}, \tag{12}$$

and $\hat{\gamma}^M\equiv V_\alpha^M\gamma^\alpha$, with $V_\alpha^M=\mathrm{diag}(e^\sigma,e^\sigma,e^\sigma,e^\sigma,1)$ the inverse vierbein. The bulk mass term M_f in Eq. (11) is expected to be of order $k\simeq M_P$. Although the fermion field Ψ is nonchiral, we can still define $\Psi_{L,R}\equiv\frac{1}{2}(1\mp\gamma_5)\Psi$. The KK decomposition can be written as

$$\Psi_{L,R}(x,y)=\frac{1}{\sqrt{2\pi r}}\sum_{n=0}\psi_n^{L,R}(x)e^{2\sigma}f_n^{L,R}(y), \tag{13}$$

where $\psi_n^{L,R}(x)$ corresponds to the nth KK fermion excitation and is a chiral four-dimensional field. Demanding that the KK fermions have the usual action in four dimensions leads to the coupled differential equations [4,5]

$$(\partial_y-M_f)f_n^L(y)=-M_n e^\sigma f_n^R(y)$$
$$(\partial_y+M_f)f_n^R(y)=M_n e^\sigma f_n^L(y), \tag{14}$$

where M_n is the mass of the nth KK fermion excitation. The corresponding normalization condition reads

$$\frac{1}{\pi r}\int_0^{\pi r}dy\, e^\sigma f_n^{L,R}(y)\, f_m^{L,R}(y)=\delta_{nm}, \tag{15}$$

where we have made use of the fact that the $f_n(y)$ have definite Z_2 parity. The zero mode wave functions are obtained from Eq. (14) for $M_n=0$. They are given by

$$f_0^{L,R}(y)=\sqrt{\frac{k\pi r\,(1\pm 2\nu_f)}{e^{k\pi r(1\pm 2\nu_f)}-1}}e^{\pm\nu_f ky}, \tag{16}$$

with $\nu_f\equiv M_f/k$ parametrizing the bulk fermion mass in units of the inverse AdS radius k. The Z_2 orbifold projection is used so that only one of these is actually allowed, either a left-handed or a right-handed zero mode.

The final piece of the bulk SM is the Higgs field. If it is allowed to live in the bulk, it gives a bulk mass term to the W and Z gauge bosons. In order to obtain the correct values of M_W and M_Z this bulk mass would have to be extremely fine tuned [4], in effect recovering the same amount of tuning as in the SM. In order to avoid this problem, the Higgs field should be localized on the TeV boundary at $y=\pi r$. The picture of the five-dimensional SM in the RS scenario is particularly attractive when we take into account its potential to generate the hierarchy of fermion masses from $\mathcal{O}(1)$ flavor breaking in the fermion bulk mass parameter ν_f. This was first considered in Ref. [6] and further examined in [7]. These authors have shown that allowing different values of ν_f within the natural constraint $|\nu_f|\simeq\mathcal{O}(1)$ results in exponentially generated flavor hierarchies. For instance, one can generate the top quark mass with [6,7] $\nu_t\simeq 0.5$, and the electron mass with $\nu_e\simeq -0.5$.

Other model building extensions, such as supersymmetry in the bulk [6], grand unification [11] and dynamical electroweak symmetry breaking [12] were considered, just to name a few. Thus, it is of great interest to study in detail what are the limitations of putting matter in the bulk.

Constraints on the bulk SM vary according to the localization of fermions in the fifth dimension. This is parametrized by the bulk mass parameter ν_f. Very large positive values of ν_f correspond to fermions highly localized on the TeV boundary at $y=\pi r$, whereas negative ν_f corresponds to fermions with larger wave functions around the $y=0$ (Planck) boundary. When fermions are localized in the TeV boundary they couple strongly to KK excitations of bulk gauge bosons, with the enhancement over the gauge coupling being given approximately by [3] $\simeq\sqrt{2\pi kr}\simeq 8.4$, resulting in a bound of $m_1>23$ TeV for the mass of the first gauge boson excitation. This constraint is obtained from a fit to electroweak observables of the SM, including the effect of KK excitations through the parameter V defined in Ref. [3] and arising from the tree-level exchange of the KK excitations. However, when fermions are allowed to be in the bulk these bounds can be greatly relaxed. This was noted in Refs. [4,6], where the bound for the first KK mode of gauge bosons is given as $m_1\gtrsim 2.1(g_1/g)$ TeV, with the ratio (g_1/g) depending on the value of the bulk fermion mass parameter ν_f. The localized fermion result is recovered for large positive ν_f, $\nu_f=0$ gives $m_1\gtrsim 9$ TeV, and for negative values the bounds are much weaker, allowing the 1 TeV mass range.[1] In the next section we will see that new non-oblique effects result in considerably stronger bounds.

III. THE EFFECTS OF NON-LOCAL WAVE-FUNCTION RENORMALIZATION

In the presence of the Higgs VEV on the TeV boundary, the "zero modes" of the W and Z gauge bosons are no longer flat in the fifth dimension [7]. The resulting localized mass term repels the wave function in the vicinity of the TeV boundary. In Ref. [10] it was found that this leads to large tree level contributions to the oblique parameters S and T. The Z wave function acquires a dependence on y given by [10]

$$\chi_Z^{(0)}(y)\simeq 1+\frac{M_Z^2}{4\Lambda_r^2}\{2\pi kr-1+(1-2ky)\,e^{2k(y-\pi r)}\}, \tag{17}$$

where we have assumed $M_Z\ll\Lambda_r$, and the expression for the W wave function is obtained by replacing $M_Z\to M_W$. This leads to contributions to S and T which are approximately [10]

$$S\simeq -4\pi\frac{v^2}{\Lambda_r^2}k\pi r, \tag{18}$$

$$T\simeq -\frac{\pi}{2c_w^2}\frac{v^2}{\Lambda_r^2}k\pi r, \tag{19}$$

where $v\simeq 246$ GeV, $c_w\equiv g/\sqrt{g^2+g'^2}$, and g,g' are the four-dimensional $SU(2)_L$ and $U(1)_Y$ gauge couplings, respectively. In addition, the KK excitations of gauge bosons induce four-fermion interactions parametrized by a shift in G_F given by [8]

$$V=\sum_{n=1}^{\infty}\left(\frac{g_f^n(\nu)}{g_f^{\rm SM}}\right)^2\frac{M_W^2}{m_n^2}, \tag{20}$$

where g_f^n denotes the coupling of the nth gauge boson excitation to zero-mode fermions, and $g_f^{\rm SM}$ the corresponding SM coupling. Since, unlike in Ref. [10], we are considering bulk fermions, their couplings g_f^n will depend on the bulk mass parameter ν. The value of V obtained in Ref. [10] is recovered in the large positive ν limit. In that limit, all KK excitations couple with $g_f^n\simeq g_f^{\rm SM}\sqrt{2\pi kr}$. However, with fermions in the bulk and $|\nu|\simeq O(1)$, only the first KK excitation couples strongly, and Eq. (20) can be approximated by its first term. This results in

$$V\simeq\frac{g^2}{12}\frac{v^2}{\Lambda_r}k\pi r I^2(\nu), \tag{21}$$

where we defined

$$I(\nu)\equiv\frac{1+2\nu}{1-e^{-k\pi r(1+2\nu)}}\times\int_0^1 u^{1+2\nu}\frac{J_1(x_1u)+\alpha_1 Y_1(x_1u)}{|J_1(x_1)+\alpha_1 Y_1(x_1)|}, \tag{22}$$

with α given in Eq. (8) and x_1 defined by Eq. (9).

In Ref. [10] the contributions from S, T and V are fit to electroweak observables, in a scenario with only gauge fields propagating in the bulk. For $m_h=115$ GeV they obtained the bound [10] $\Lambda_r>11$ TeV at 95% confidence level (C.L.). This constraint on the low energy scale Λ_r translates into a lower bound on the lightest gauge boson KK excitation of $m_1>27$ TeV, and is slightly stronger than the ones previously obtained for fermions on the TeV brane and where only the V parameter had been considered.

Here we show that when fermions are allowed in the bulk there are additional non-oblique effects in the couplings of W and Z to fermions. The electroweak gauge boson wave functions are normalized to their SM values at the low energy boundary. However, the couplings to fermions are a non-local quantity resulting from the overlap of the gauge boson and fermion wave functions in the bulk. This effect is present even in the couplings to zero mode fermions, and is due to the y dependence of the gauge boson wave functions. In general the $5D$ coupling of fermions to a bulk gauge boson $\mathcal{G}$ is given by

$$S_{\rm int}=g_5\int d^4x\,dy\,\sqrt{-g}\,\mathcal{G}_\mu(x,y)\times\bar{\Psi}_f(x,y)\,\hat{\gamma}^\mu\Psi_f(x,y), \tag{23}$$

where $g_5=\sqrt{2\pi r}$, and f is a fermion flavor index. This index has been kept since in principle it is possible that the fermion wave functions (16) are flavor dependent if the bulk mass parameter ν_f is not universal. The coupling of the gauge boson to a given fermion f relative to its SM value is then given by

[1] It should be noted that these bounds are on the mass of the first KK excitation of a gauge boson. But since $m_1\simeq 2.45\Lambda_r$ (i.e. for $e^{k\pi r}=10^{16}$), the corresponding bounds on Λ_r are weaker.

$$\left(\frac{g_f}{g_f^{\mathrm{SM}}}\right)=\frac{1}{\pi r}\int_0^{\pi r}dy\, e^{ky}\,|f_0^A(y)|^2\,\chi_G^{(0)}(y), \tag{24}$$

with $A=(L,R)$ and $G=(W,Z)$. We define the new parameter

$$\gamma_f^G\equiv\left(\frac{g_f}{g_f^{\mathrm{SM}}}\right)-1, \tag{25}$$

where $f=b_L,b_R,t_L,\ldots$ is the zero mode fermion label. From Eqs. (16) and (17) we have

$$\gamma_f^Z=\frac{|f_0^A(0)|^2}{4k\pi r}\left(\frac{M_Z}{\Lambda_r}\right)^2(I_1+I_2+I_3), \tag{26}$$

where we have defined

$$I_1=\frac{(2k\pi r-1)}{1\pm 2\nu_f}[e^{(1\pm 2\nu_f)k\pi r}-1],$$

$$I_2=\frac{e^{-2k\pi r}}{3\pm 2\nu_f}[e^{(3\pm 2\nu_f)k\pi r}-1],$$

$$I_3=-\frac{2e^{-2k\pi r}}{3\pm 2\nu_f}\left\{\left(k\pi r-\frac{1}{3\pm 2\nu_f}\right)e^{(3+2\nu_F)k\pi r}+\frac{1}{3\pm\nu_f}\right\}, \tag{27}$$

and the $+(-)$ corresponds to left-handed (right-handed) fermions. A similar expression can be obtained for the shift in the W couplings by noting that $\gamma_f^W=c_w^2\gamma_f^Z$. We first note that γ_f^G always defines a *positive* shift in the corresponding coupling. We can also see that for $\nu_f<-0.5$, the quantity γ_f^G is practically independent of ν_f and is given by

$$\gamma_f^Z\simeq\frac{M_Z^2}{4\Lambda_r^2}(2\pi kr-1)\qquad(\nu_f<-0.5). \tag{28}$$

For larger values of ν_f, the value of γ_f^Z is reduced by the ν_f-dependent terms.

In order to study the effects of γ_f^Z on the bounds on the scale Λ_r, we will first assume that the bulk fermion mass is universal. In the next section we will study the effects of flavor violation in the bulk and the possible signals associated with it. We then consider, for the remaining of this section, the case where $\nu_{f_L}=-\nu_{f_R}\equiv\nu$ for all fermions propagating in the bulk. With this choice, all couplings undergo a universal shift given by

$$g_f\rightarrow g_f(1+\gamma^G), \tag{29}$$

where G refers to either the W or the Z coupling. Thus, there is a universal shift in the charged current couplings such that

$$\frac{G_F}{\sqrt{2}}=\frac{e^2}{8s_w^2c_w^2M_Z^2}(1+\gamma^W)^2. \tag{30}$$

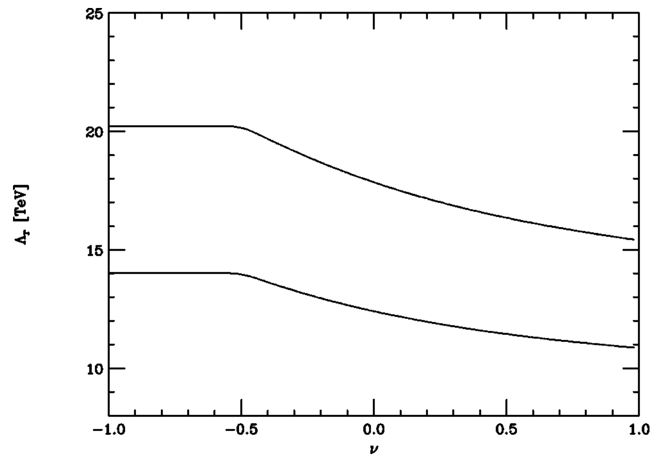

FIG. 1. Lower bounds (95% C.L.) on Λ_r vs the fermion bulk mass parameter ν. The top curve corresponds to $m_h=115$ GeV, the lower curve is for $m_h=300$ GeV.

Following the standard procedure [13,14], we redefine the Weinberg angle by

$$s_wc_w(1-\gamma^W)\rightarrow s_wc_w. \tag{31}$$

This means that we recover the standard form of the relation between G_F, α, s_w and M_Z, whereas the neutral current coupling now is

$$\frac{e}{s_wc_w}\rightarrow\frac{e}{s_wc_w}(1+\gamma^Z-\gamma^W)=\frac{e}{s_wc_w}(1+s_w^2\gamma^Z). \tag{32}$$

The replacement in Eq. (31) implies that there is a new contribution to s_w^2 given by

$$s_w^2\rightarrow s_w^2\left(1+\cdots+\frac{c_w^2}{c_w^2-s_w^2}2\gamma^W\right), \tag{33}$$

where the dots stand for the contributions from the S, T and V parameters. Equation (33) implies an additional shift in fermion couplings to the Z through the factor $(T_3^f-Q_fs_w^2)$.

We perform a fit to the electroweak observables listed in the Appendix, where we also show the dependence on S, T, V and γ^Z. The data are taken from Ref. [15]. For a fixed value of the fermion bulk mass parameter ν we obtain bounds on the low energy scale Λ_r. In Fig. 1 we plot the 95% C.L. bound on Λ_r as a function of ν. The top curve corresponds to $m_h=115$ GeV. We can see that the addition of the parameter γ^Z arising from fermion delocalization results in stronger bounds. The constraint obtained in Ref. [10] for fermions localized on the low energy brane ($\Lambda_r>11$ TeV), is recovered in the $\nu\gg 1$ limit. It was pointed out in Refs. [6,16] that the hierarchy of fermion masses could be naturally obtained in the bulk SM for values of $\nu_f<-0.5$ for all fermions except the top quark. Since the flavor dependence of V and γ^Z is exponentially suppressed for these values, the bounds in Fig. 1 apply to this model. Thus the 95% C.L. limit on Λ_r in this scenario for generating fermion masses is $\Lambda_r>20$ TeV, which translates into a mass bound for the first KK excitation of the gauge bosons which is $m_1>49$ TeV.

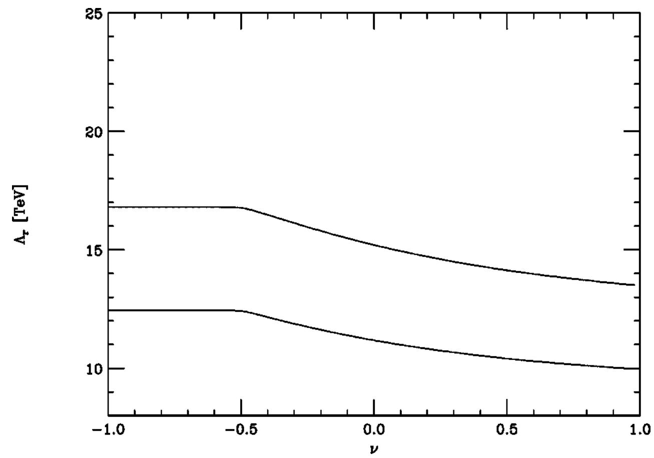

FIG. 2. Lower bounds (95% C.L.) on Λ_r vs the fermion bulk mass parameter ν for the case with the third generation confined to the low energy boundary. The top curve corresponds to $m_h = 115$ GeV, the lower curve is for $m_h = 300$ GeV.

Increasing the mass of the Higgs weakens the bounds on Λ_r somehow, as it can be seen from the bottom curve in Fig. 1, for $m_h = 300$ GeV. However, the quality of the fit worsens considerably as m_h grows. The value of the χ^2 is practically insensitive to the ν parameter. Varying then Λ_r and m_h, the minimum χ^2 is obtained for the lighter m_h and $\Delta\chi^2 \approx 6.2$ for $m_h = 300$ GeV. We then conclude that $m_h < 300$ GeV at 95% C.L. in the bulk SM of the Randall-Sundrum scenario.

We also study a scenario with the third generation localized on the low energy boundary. It was recently suggested in Ref. [9] that this is needed in order to avoid potentially large contributions to the T parameter from the KK excitations of the top quark. In Fig. 2 we plot the 95% C.L. bounds on Λ_r, where the fermion bulk mass parameter ν now refers to that of the first two generations which live in the bulk. Since the third generation does not propagate in the bulk, it does not feel the effects of γ^Z. The resulting bounds are somewhat lower than the ones in Fig. 1. However, they are roughly three times stronger than the ones derived in [9]. For instance, for $\nu < -0.5$ Fig. 2 implies that the first gauge boson KK excitation must obey $m_1 > 41$ TeV at 95% C.L. and for $m_h = 115$ GeV. The lower curve corresponds to $m_h = 300$ GeV and, just as for the case of Fig. 1, corresponds to $\Delta\chi^2 \simeq 6.2$ and thus represents the 95% C.L. value for m_h in a fit of Λ_r and m_h.

We end this section with a comment on the possible effects of higher dimensional operators. In principle, since the 5D theory is non-renormalizable we can write down higher-dimensional operators suppressed by the appropriate powers of the relevant 5D scale. As an example we consider the operator

$$\frac{c_1}{M_5}[\bar{\Psi}_f(x,y)\hat{\sigma}^{MN}\Psi_f(x,y)]\mathcal{G}_{MN}(x,y) \tag{34}$$

where $M_5 \simeq M_P$ is the cut-off scale of the effective 5D theory and c_1 is a dimensionless coefficient which is naturally $c_1 \simeq \mathcal{O}(1)$. Although this is suppressed with respect to the leading operator in Eq. (23), it is possible that in reducing to the 4D effective theory the resulting operator may be suppressed only by the TeV scale due to the presence of the warp factor. For instance, projecting onto the zero modes results in the interaction

$$\frac{c_1}{M_5}\frac{1}{\sqrt{2\pi M_5 r}}\frac{\left(\frac{1}{2}+\nu\right)}{2(1+\nu)}\frac{e^{2k\pi r(1+\nu)}-1}{e^{k\pi r(1+2\nu)}} \times(\bar{f}_R\sigma^{\mu\nu}f_L)Z_{\mu\nu}. \tag{35}$$

This implies a contribution to the Z couplings to the fermion f, which for the values of ν considered here gives approximately

$$\frac{m_f}{M_Z}\ \text{few}\times 10^{-4}c_1, \tag{36}$$

where we considered on-shell fermions. As usual, we impose $k < M_5/\sqrt{20}$ (resulting from asking that the curvature be smaller than M_5^2) so there are no effects due to strong gravitational interactions. Then, for the couplings we consider here this results in shifts that are typically $\lesssim \text{few}\times 10^{-6}c_1$. This is at least two orders of magnitude smaller than the effects implied by the values of γ^Z in Eqs. (26) and (28). We conclude that this particular operator can be safely ignored. Nonetheless, we see that it is not suppressed by the Planck scale relative to the leading operator in Eq. (23). This exercise highlights the fact that the effects of higher dimensional operators are not *a priori* to be ignored and that, in some cases, they could have important effects at the weak scale. However, for the purpose of our analysis, we assume these will not change considerably the constraints derived in this section.

IV. EFFECTS OF FLAVOR VIOLATION IN THE BULK

As was mentioned in the previous section, it is possible to generate a large hierarchy in fermion masses if we allow $\mathcal{O}(1)$ flavor breaking in the bulk. Then we may allow the bulk fermion mass parameters ν_f to vary as long as they all are of order one (i.e. all bulk fermions have masses of the order of M_P). This means that the shift in fermion couplings to the W and Z gauge bosons given in Eq. (26) may be non-universal. This necessarily leads to flavor-changing neutral currents (FCNC) of the Z at tree level, as well as non-universal corrections to the charged current interactions. We concentrate here on the FCNC of the Z due to their dangerous phenomenological nature. We first show how FCNC come about in the present context. Assuming that γ_f^Z is different for each fermion induces a flavor-dependent shift in the Z coupling given by

$$\delta g_f = \gamma_f \frac{e}{s_w c_w}(T_3^f - Q_f s_w^2), \tag{37}$$

where T_3^f is the third component of weak isospin, Q_f is the fermion electric charge and we have dropped the superscript

Z in γ_f. The non-universality in these shifts results in FCNC when fermions are rotated from the weak to the mass eigenbasis.

Let us first consider the down quark sector. In the weak eigenbasis, the Z coupling to down quark types reads

$$\mathcal{L}=-\frac{e}{s_w c_w}Z^{\mu}\left\{\left(-\frac{1}{2}+\frac{s_w^2}{3}\right)\times\sum_{D=d,s,b}(1+\gamma_{D_L})(\overline{D}_L\gamma_\mu D_L)+\frac{s_w^2}{3}\sum_{D=d,s,b}(1+\gamma_{D_R})\times(\overline{D}_R\gamma_\mu D_R)\right\}. \quad (38)$$

In principle, in Eq. (38) there is also a factor of $(1-\gamma^W)$ coming from the effect in charged currents, as in Eq. (32). However, this will be a flavor universal shift (as far as the quark flavor goes[2]). As we will see below, this will cancel in the FCNC effects since these will depend on the differences of quark flavor-dependent quantities. We define the rotation of down quarks into the mass eigenbasis by $D_L\to A^L D_L$ and $D_R\to A^R D_R$ [here $D^T\equiv(dsb)$]. There will be analogous rotation matrices in the up sector given by $U_L\to B^L U_L$, etc., such that $V_{CKM}=(B^L)^\dagger A^L$ is the usual quark mixing matrix. The unitarity of A^L and A^R implies that flavor off-diagonal terms not proportional to a factor of γ_D vanish. Then, off-diagonal terms are given by

$$\mathcal{L}_D^{\rm FCNC}=-\frac{e}{s_w c_w}Z^\mu\{\Delta_L^{ds}(\overline{d}_L\gamma_\mu s_L)+\Delta_L^{db}(\overline{d}_L\gamma_\mu b_L)+\Delta_L^{sb}(\overline{s}_L\gamma_\mu b_L)+(\mathrm{L}\to\mathrm{R})+\mathrm{H.c.}\}, \quad (39)$$

where we defined

$$\Delta_L^{ds}\equiv\gamma_{d_L}A_{11}^{L*}A_{12}^L+\gamma_{s_L}A_{21}^{L*}A_{22}^L+\gamma_{b_L}A_{31}^{L*}A_{32}^L, \quad (40)$$

$$\Delta_L^{db}\equiv\gamma_{d_L}A_{11}^{L*}A_{13}^L+\gamma_{s_L}A_{21}^{L*}A_{23}^L+\gamma_{b_L}A_{31}^{L*}A_{33}^L, \quad (41)$$

$$\Delta_L^{sb}\equiv\gamma_{d_L}A_{12}^{L*}A_{13}^L+\gamma_{s_L}A_{22}^{L*}A_{23}^L+\gamma_{b_L}A_{32}^{L*}A_{33}^L. \quad (42)$$

The analogous expressions for Δ_R^{ij} can be obtained by $L\to R$ in Eqs. (40)–(42). Once again, we notice that if $\gamma_d=\gamma_s=\gamma_b$, the unitarity of $A^{L,R}$ implies that all off-diagonal terms vanish. Unitarity also implies that Eqs. (40)–(42) actually depend on two independent combinations of γ_Ds, e.g. $(\gamma_{d_L}-\gamma_{s_L})$ and $(\gamma_{b_L}-\gamma_{s_L})$. Then, as mentioned above, any universal shift in the Z couplings (38) cancels out when considering the flavor changing terms, Eq. (39). The interactions in Eqs. (40)–(42) will induce K^0-$\overline{K}^0$, B_d^0-$\overline{B}_d^0$ and B_s^0-$\overline{B}_s^0$ mixing, as well as rare K and B decays, all mediated by tree-level Z exchange. Similar expressions can be written for the up quark sector.

In principle, we have little information on the entries of $A^{L,R}$ or $B^{L,R}$. In order to illustrate how FCNC are generated let us examine a simple model for the rotation matrices. Let us consider the situation where $B^L\simeq I$ and $A^L\simeq V_{\rm CKM}$. If this is the case then we have, for instance for the $d_L\to s_L$ term,

$$\Delta_L^{ds}\simeq\gamma_{d_L}V_{ud}^*V_{us}+\gamma_{s_L}V_{cd}^*V_{cs}+\gamma_{b_L}V_{td}^*V_{ts}\simeq\lambda(\gamma_{d_L}-\gamma_{s_L}), \quad (43)$$

where the last line results from $V_{ud}^*V_{us}+V_{cd}^*V_{cs}\simeq 0$ and $\lambda\simeq 0.22$ is the sine of the Cabibbo angle. We would also obtain

$$\Delta_L^{db}\simeq(\gamma_{d_L}-\gamma_{s_L})V_{ud}^*V_{ub}+(\gamma_{b_L}-\gamma_{s_L})V_{td}^*V_{tb}$$

$$\Delta_L^{sb}\simeq(\gamma_{d_L}-\gamma_{s_L})V_{us}^*V_{ub}+(\gamma_{b_L}-\gamma_{s_L})V_{ts}^*V_{tb}. \quad (44)$$

Although in general there is no reason to believe that $B^L\simeq I$, we will refer to this scenario (as well as to the general case) in order to evaluate the potential size of the effects in flavor physics.

We first consider FCNC processes in kaon physics induced by the Z exchange. The contribution to K^0-$\overline{K}^0$ mixing is given by

$$\Delta m_K=\frac{4G_F}{\sqrt{2}}f_K^2 m_K\left\{\frac{2}{3}\left(-\frac{1}{2}\frac{s_w^2}{3}\right)^2\mathrm{Re}[(\Delta_L^{ds})^2]+\frac{2}{3}\left(\frac{s_w^2}{3}\right)^2\mathrm{Re}[(\Delta_R^{ds})^2]-4\left(-\frac{1}{2}+\frac{s_w^2}{3}\right)\times\left(\frac{s_w^2}{3}\right)\left[\frac{1}{4}+\frac{1}{6}\left(\frac{m_K}{m_s+m_d}\right)^2\right]\mathrm{Re}[\Delta_L^{ds}\Delta_R^{ds}]\right\}. \quad (45)$$

Even in the presence of the last term, which is "chirally enhanced" and dominates, we do not find a large effect. If we assume that the Z exchange contribution saturates the experimental measurement [15] $\Delta m_K^{\rm exp}=(3.489\pm0.008)\times10^{-15}$ GeV, then considering $\Delta_L^{ds}\simeq\Delta_R^{ds}$, we find $\sqrt{\mathrm{Re}[(\Delta_L^{ds})^2]}<1\times10^{-4}$.

A tighter bound is obtained from $K_L\to\mu^+\mu^-$. This is extracted from [14]

[2]Here we actually have $(1-\frac{1}{2}\gamma_e^W-\frac{1}{2}\gamma_\mu^W)$ as entering in μ decay.

$$\frac{\mathrm{Br}(K_L\to\mu^+\mu^-)}{\mathrm{Br}(K^+\to\mu^+\nu_\mu)} = \frac{\tau(K_L)}{\tau(K^+)}\frac{8}{|V_{us}|^2}\left[\left(-\frac{1}{2}+s_w^2\right)^2+(s_w^2)^2\right] \times\left[\left(-\frac{1}{2}+\frac{s_w^2}{3}\right)^2|\Delta_L^{ds}|^2+(s_w^2)^2|\Delta_R^{ds}|^2\right]. \tag{46}$$

With [15] $\mathrm{Br}(K_L\to\mu^+\mu^-)=(7.18\pm0.17)\times10^{-9}$ and assuming that the new contribution saturates the rate, we obtain $\Delta_L^{ds}<4.8\times10^{-5}$.

But the best bound on $\Delta_{L,R}^{ds}$ comes from $K^+\to\pi^+\nu\bar{\nu}$. The contribution to the amplitude is given by

$$\delta\mathcal{A}=\frac{e^2}{s_w^2c_w^2M_Z^2}\frac{1}{2}\left\{\left(-\frac{1}{2}+\frac{s_w^2}{3}\right)\Delta_L^{ds}(\bar{d}_L\gamma_\mu s_L) +\frac{s_w^2}{3}\Delta_R^{ds}(\bar{d}_R\gamma_\mu s_R)\right\}\times\sum_{i=e,\mu,\tau}(\bar{\nu}_L^i\gamma^\mu\nu_L^i)+\mathrm{H.c.} \tag{47}$$

On the other hand, the SM amplitude can be written as [17]

$$\mathcal{A}_{\mathrm{SM}}=\frac{G_F}{\sqrt{2}}\frac{\alpha}{\pi s_w^2}S(\bar{d}_L\gamma_\mu s_L)\times\sum_{i=e,\mu,\tau}(\bar{\nu}_L^i\gamma^\mu\nu_L^i)+\mathrm{H.c.}, \tag{48}$$

where $S\simeq2\times10^{-3}$. The current experimental measurement is [18] $\mathrm{Br}(K^+\to\pi^+\nu\bar{\nu})=(1.57^{+1.75}_{-0.82})\times10^{-10}$, or an upper bound of $\mathrm{Br}(K^+\to\pi^+\nu\bar{\nu})<5\times10^{-10}$. The left-handed term in Eq. (47) interferes with the SM amplitude. If we consider this term only, we derive the 2σ bound

$$\Delta_L^{ds}<1.2\times10^{-5}. \tag{49}$$

Conversely, if we only consider the right-handed contribution we obtain

$$\Delta_R^{ds}<1.9\times10^{-5}. \tag{50}$$

Finally, if we consider $\Delta_L^{ds}\simeq\Delta_R^{ds}$, then the bound is

$$\Delta_{L,R}^{ds}<7.5\times10^{-6}. \tag{51}$$

In order to estimate the compatibility of these bounds with the ones derived in the previous section from electroweak precision measurements, we assume $B\simeq I$ so we can make use of Eq. (43) which, together with the bound in Eq. (49), gives

$$(\gamma_{d_L}-\gamma_{s_L})<5.5\times10^{-5}. \tag{52}$$

If we now consider as reference the bounds on Λ_r that we obtained in the previous section and remember that for $\nu_f<-0.5$ γ_f is ν_f independent, we can derive bounds on bulk mass differences. For instance, for $\Lambda_r>20$ TeV, Eq. (52) implies $(\nu_d-\nu_s)<0.1$, as long as one of them is greater than -0.5. For $\Lambda_r>13$ TeV, $(\nu_d-\nu_s)<0.08$, for $\Lambda_r>10$ TeV, $(\nu_d-\nu_s)<0.05$, etc. Then, although the bounds imply a certain amount of tuning between the bulk mass parameters, this is not particularly worrisome. If both bulk mass parameters are $(\nu_d,\nu_s)<-0.5$, the bound (52) is easily satisfied. In the most general case, there is too much freedom in order to use Eqs. (49) and (50) to constraint Λ_r and the bulk mass parameters. This is particularly true in the case of Δ_R^{ds}, since—even assuming that $A^L\simeq V_{\mathrm{CKM}}$ holds—we have no information on A^R. However, the lesson we draw from the bound in Eq. (52) is that $K^+\to\pi^+\nu\bar{\nu}$ is sensitive to small flavor breaking in the bulk, particularly taking into account future improvements in the experimental measurements of this mode as well as $K_L\to\pi^0\nu\bar{\nu}$.

Observables in B physics turn out to be less sensitive. For instance, following the derivation of Eq. (45), B^0-$\bar{B}^0$ mixing results in the bound $\Delta_{L,R}^{db}<5\times10^{-4}$, where again we assumed $\Delta_L^{db}\simeq\Delta_R^{db}$. This in general corresponds to a much weaker constraint on $(\gamma_{d_L}-\gamma_{b_L})$, as it is illustrated by the scenario with $A^L\simeq V_{\mathrm{CKM}}$. In this case, $\Delta_L^{db}\simeq\lambda^3(\gamma_{d_L}-\gamma_{b_L})$, so that the effect is largely suppressed by the large power of the Cabibbo angle. The bound gets weaker if we assume only one of them non-vanishing, since the last term in Eq. (45) is still very important (although not dominant). Although Δ_R^{db} could be large and unsuppressed by powers of λ, its effect in B^0-$\bar{B}^0$ mixing is suppressed—relative to that of Δ_L^{sb}—by two powers of $(s_w^2/3)/(-1/2+s_w^2)\simeq0.18$. The current bounds from rare semileptonic B decays such as $b\to sl^+l^-$ are even weaker. Although a great deal of improvement in the experimental situation is expected soon in these decay modes, these are still not sensitive to $\Delta_{L,R}^{sb}$ once we consider the bounds on Λ_r from the previous section.

Finally, we consider the effects of flavor violation in the up quark sector. There is an expression analogous to Eq. (38) for the Z couplings to up quarks, which can be obtained by replacing D by $U=(u\ c\ t)^T$, and by putting the appropriate factors of $(T_3-Q_fs_w^2)$. There are also expressions for $\Delta_{L,R}^{uc}$, $\Delta_{L,R}^{ut}$ and $\Delta_{L,R}^{ct}$ similar to those in Eqs. (40)–(42), where the rotation matrix is now $B^{L,R}$. The terms involving the top quark will only surface in top rare decays, which require top data samples not yet available. We will then concentrate on the uc terms. Current experimental constraints on D^0-$\bar{D}^0$ mixing [19] result in[3] $\Delta m_D<4.5\times10^{-14}$ GeV. The contributions of $\Delta_{L,R}^{uc}$ to Δm_D can be read off Eq. (45). If we assume $\Delta_L^{uc}\simeq\Delta_R^{uc}$ we obtain

$$\Delta_{L,R}^{uc}<3\times10^{-4}. \tag{53}$$

Unlike in the case of $\Delta_{L,R}^{db}$, here we do not expect a large suppression by powers of λ. For instance, if we assume that all of the CKM matrix comes from the up sector, i.e. B^L

[3]This bound is obtained assuming that there is no strong relative phase between the doubly Cabibbo suppressed decay $D^0\to K^+\pi^-$ and the Cabibbo allowed mode $D^0\to K^-\pi^+$.

$\simeq V_{\mathrm{CKM}}$ and $A^L \simeq I$, then $\Delta_L^{uc} \simeq \lambda(\gamma_{u_L} - \gamma_{c_L})$, only suppressed by λ. Moreover, it is possible to imagine that Δ_R^{uc} is unsuppressed, which would mean that D^0-$\bar{D}^0$ mixing is sensitive to $(\gamma_{u_L} - \gamma_{c_L}) \simeq 3\times 10^{-4}$. If we assume that only Δ_L^{uc} is non-vanishing, we obtain $\Delta_L^{uc} < 5\times 10^{-4}$. If, on the other hand, we assume that only the right-handed term is present, we obtain $\Delta_R^{uc} < 5\times 10^{-3}$. From the bounds derived in the previous section we know that $\gamma_f \lesssim 1\times 10^{-3}$ for $\nu_f < -0.5$, and that the flavor dependence reduces γ_f as ν_f becomes more positive. We then conclude that the flavor universal bounds from electroweak precision measurements are compatible with the bounds on $\Delta_{L,R}^{uc}$ from Δm_D, even for relatively large values of $(\nu_u - \nu_c)$, the difference in the relevant bulk mass parameters. Future improvements in the experimental bound on Δm_D will begin to probe flavor violation in the bulk in a region where the flavor dependence in γ_f could manifest itself. However, at the moment this is not a constraint comparable to the ones derived above from kaon processes. On the other hand, rare D decays receive very small contributions from $\Delta_{L,R}^{uc}$ once the bounds derived from Δm_D are taken into account. For instance, in $D \to (\pi,\rho) l^+ l^-$, they mainly contribute to the four-fermion operator $(\bar{u}_L \gamma_\mu c_L) \times (\bar{l} \gamma^\mu \gamma_5 l)$. Even when looking at the low dilepton mass region, away from the dominating resonant contributions, the effects of $\Delta_{L,R}^{uc}$ are small enough to be comparable to the remaining hadronic uncertainties present [20].

V. CONCLUSIONS

We have derived constraints on the low energy scale Λ_r induced in the RS scenario, when the SM fields are allowed to propagate in the 5D bulk. The effects are the result of the deformation of the Z and W zero-mode wave functions due to the presence of the Higgs field in the low energy boundary. When only gauge fields are allowed in the bulk, this only leads to the contributions to the oblique parameters S and T discussed in Ref. [10]. If fermions also propagate in the bulk, non-oblique effects arise due to the modified couplings of the Z and the W to *zero mode* fermions. We considered these new effects here. In Sec. III we studied the constraints one obtains by assuming that the shifts in the couplings to fermions are flavor universal. This is a rather good assumption, as we point out in the discussion leading to Eq. (28): for values of the fermion bulk mass parameter $\nu_f < -0.5$, corrections to it are exponentially suppressed. This is the region of ν_f that can be used in order to generate the fermion masses as pointed out in Refs. [6,16]. For $\nu_f > -0.5$, the dependence of the bounds on ν_f is rather mild. Since the mass of the first KK excitation of a gauge boson is given by $m_1 \simeq 2.45\Lambda_r$, we conclude that these must be heavier than a few tens of TeV, putting them out of reach of the Large Hadron Collider at CERN.

We also found—just as in Ref. [10] for the case of localized fermions—that although the bounds on Λ_r are somewhat relaxed as the Higgs mass increases, the χ^2 considerably worsens. We observe in the bulk SM $m_h < 300$ GeV at 95% C.L., a bound very similar to the one derived in [10].

We also considered the scenario motivated in Ref. [9], where the third generation is confined to the low energy boundary. The bounds on Λ_r are displayed in Fig. 2. Although these are somehow lower than those in Fig. 1, they are still about three times stronger than the ones obtained in Ref. [9].

In Sec. IV, we considered the residual effects of flavor dependence in γ_f^Z by allowing $\mathcal{O}(1)$ flavor breaking in the fermion bulk mass parameters ν_f. This induces FCNC of the Z to zero-mode fermions. We showed that these FCNC effects are not dangerous once the constraints on Λ_r derived in Sec. III are taken into account. We found the most sensitive observables to be in kaon physics such as K^0-$\bar{K}^0$ mixing, $K_L \to \mu^+\mu^-$, but especially $K^+ \to \pi^+ \nu\bar{\nu}$. The constraints from these imply that there should be a certain amount of degeneracy in bulk masses of the down quark sector. Although the bounds do not result in fine tuning, better measurements of $K^+ \to \pi^+ \nu\bar{\nu}$ may imply the need of a higher Λ_r if one is to avoid 1% adjustments in the bulk masses.

We stress that the effects considered here occur at tree level and are the result of interactions among zero-mode gauge bosons and fermions. Additional contributions to oblique and non-oblique parameters may result from loop effects involving KK excitations. However, in most cases we do not expect the predictions to be well defined. In particular, we expect that cut off dependence would hinder our ability to translate a one loop calculation into a bound on the parameters of the theory. This sensitivity to the cut-off can be interpreted as a consequence of the fact that the 5D theory is non-renormalizable, so higher-dimensional operators with unknown coefficients may absorb the cut-off dependence from a naive loop computation. As we show at the end of Sec. III, the effect of higher-dimensional operators could partially cancel some of the effects discussed here. However, it appears unnatural to expect that they could do so efficiently enough to loosen the constraint considerably.

There are also flavor-violating effects induced by the mixing of zero-mode fermions with KK excitations [21]. However, these become irrelevant once Λ_r is raised to the values shown in Figs. 1 and 2.

ACKNOWLEDGMENTS

The author thanks Mike Chanowitz and John Terning for useful comments, and especially Walter Goldberger for useful discussions and suggestions during the course of this work. This work was supported by the Director, Office of Science, Office of High Energy and Nuclear Physics of the U.S. Department of Energy under Contract DE-AC0376SF00098.

APPENDIX

What follows are the expressions of observables used in the fits of Sec. III as functions of S, T, V and γ^Z:

$$\Gamma_Z = (\Gamma_Z)^{\mathrm{SM}}(1 - 3.8\times 10^{-3} S + 0.011 T - 1.4 V - 0.08\gamma^Z)$$

$$R_e=(R_e)^{\mathrm{SM}}(1-2.9\times10^{-3}S+2\times10^{-3}T-0.26V-0.4\gamma^Z)$$

$$R_\mu=(R_\mu)^{\mathrm{SM}}(1-2.9\times10^{-3}S+2\times10^{-3}T-0.26V-0.4\gamma^Z)$$

$$R_\tau=(R_\tau)^{\mathrm{SM}}(1-2.9\times10^{-3}S+2\times10^{-3}T-0.26V-0.4\gamma^Z)$$

$$\sigma_h=(\sigma_h)^{\mathrm{SM}}(1+2.2\times10^{-4}S-1.6\times10^{-4}T-0.021V-0.96\gamma^Z)$$

$$\Gamma_b=(\Gamma_b)^{\mathrm{SM}}(1-4.5\times10^{-3}S+0.011T-1.4V-0.18\gamma^Z)$$

$$\Gamma_c=(\Gamma_c)^{\mathrm{SM}}(1-6.5\times10^{-3}S+0.0124T-1.6V+0.45\gamma^Z)$$

$$\Gamma_{\mathrm{inv.}}=(\Gamma_{\mathrm{inv.}})^{\mathrm{SM}}(1+7.8\times10^{-3}T-V+0.46\gamma^Z)$$

$$A_{\mathrm{FB}}^e=(A_{\mathrm{FB}}^e)^{\mathrm{SM}}-6.8\times10^{-3}S+4.8\times10^{-3}T-0.62V-0.95\gamma^Z$$

$$A_{\mathrm{FB}}^\mu=(A_{\mathrm{FB}}^\mu)^{\mathrm{SM}}-6.8\times10^{-3}S+4.8\times10^{-3}T-0.62V-0.95\gamma^Z$$

$$A_{\mathrm{FB}}^\tau=(A_{\mathrm{FB}}^\tau)^{\mathrm{SM}}-6.8\times10^{-3}S+4.8\times10^{-3}T-0.62V-0.95\gamma^Z$$

$$A_\tau(P_\tau)=(A_\tau(P_\tau))^{\mathrm{SM}}-0.028S+0.020T-2.6V-4\gamma^Z$$

$$A_e(P_\tau)=(A_e(P_\tau))^{\mathrm{SM}}-0.028S+0.020T-2.6V-4\gamma^Z$$

$$A_{\mathrm{FB}}^b=(A_{\mathrm{FB}}^b)^{\mathrm{SM}}-0.020S+0.014T-1.8V-2.77\gamma^Z$$

$$A_{\mathrm{FB}}^c=(A_{\mathrm{FB}}^c)^{\mathrm{SM}}-0.016S+0.011T-1.4V-2.16\gamma^Z$$

$$A_{\mathrm{LR}}=(A_{\mathrm{LR}})^{\mathrm{SM}}-0.028S+0.02T-2.6V-4\gamma^z$$

$$M_W=(M_W)^{\mathrm{SM}}(1-3.6\times10^{-3}S+5.5\times10^{-3}T-0.71V+0.66\gamma^Z)$$

$$g_L^2(\nu N\to\nu N)=[g_L^2(\nu N\to\nu N)]^{\mathrm{SM}}-2.7\times10^{-3}S+6.5\times10^{-3}T-0.066V-0.096\gamma^Z$$

$$g_R^2(\nu N\to\nu N)=[g_R^2(\nu N\to\nu N)]^{\mathrm{SM}}+9.3\times10^{-4}S+2.0\times10^{-4}T+0.1V+0.16\gamma^Z$$

$$g_{eV}(\nu e\to\nu e)=[g_{eV}(\nu e\to\nu e)]^{\mathrm{SM}}+7.2\times10^{-3}S-5.4\times10^{-3}T+0.65V+0.99\gamma^Z$$

$$g_{eA}(\nu e\to\nu e)=[g_{eA}(\nu e\to\nu e)]^{\mathrm{SM}}-3.9\times10^{-3}T+0.15V-0.23\gamma^Z$$

$$Q_W(Cs)=[Q_W(Cs)]^{\mathrm{SM}}-0.793S-0.009T-0.47V-290.8\gamma^Z.$$

[1] N. Arkani-Hamed, S. Dimopoulos, and G.R. Dvali, Phys. Lett. B **429**, 263 (1998); I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos, and G.R. Dvali, *ibid.* **436**, 257 (1998).
[2] L. Randall and R. Sundrum, Phys. Rev. Lett. **83**, 3370 (1999).
[3] W.D. Goldberger and M.B. Wise, Phys. Rev. D **60**, 107505 (1999); H. Davoudiasl, J.L. Hewett, and T.G. Rizzo, Phys. Lett. B **473**, 43 (2000); A. Pomarol, *ibid.* **486**, 153 (2000).
[4] S. Chang, J. Hisano, H. Nakano, N. Okada, and M. Yamaguchi, Phys. Rev. D **62**, 084025 (2000).
[5] Y. Grossman and M. Neubert, Phys. Lett. B **474**, 361 (2000).
[6] T. Gherghetta and A. Pomarol, Nucl. Phys. **B586**, 141 (2000).
[7] S.J. Huber and Q. Shafi, Phys. Rev. D **63**, 045010 (2001).
[8] H. Davoudiasl, J.L. Hewett, and T.G. Rizzo, Phys. Rev. D **63**, 075004 (2001).
[9] J.L. Hewett, F.J. Petriello, and T.G. Rizzo, "Precision measurements and fermion geography in the Randall-Sundrum model revisited," hep-ph/0203091.
[10] C. Csáki, J. Erlich, and J. Terning, Phys. Rev. D **66**, 064021 (2002).
[11] A. Pomarol, Phys. Rev. Lett. **85**, 4004 (2000); L. Randall and M.D. Schwartz, J. High Energy Phys. **11**, 003 (2001); Phys. Rev. Lett. **88**, 081801 (2002); W.D. Goldberger and I.Z. Rothstein, "High energy field theory in truncated AdS backgrounds," hep-th/0204160.
[12] N. Rius and V. Sanz, Phys. Rev. D **64**, 075006 (2001).
[13] M.E. Peskin and T. Takeuchi, Phys. Rev. Lett. **65**, 964 (1990).
[14] C.P. Burgess, S. Godfrey, H. Konig, D. London, and I. Maksymyk, Phys. Rev. D **49**, 6115 (1994).
[15] Particle Data Group, D.E. Groom *et al.*, Eur. Phys. J. C **15**, 1 (2000); and updates from http://www-pdg.lbl.gov
[16] S.J. Huber and Q. Shafi, Phys. Lett. B **498**, 256 (2001).
[17] G. Buchalla, A.J. Buras, and M.E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996).
[18] E787 Collaboration, S. Adler *et al.*, Phys. Rev. Lett. **88**, 041803 (2002); Phys. Lett. B **537**, 211 (2002).
[19] CLEO Collaboration, G. Brandenburg *et al.*, Phys. Rev. Lett. **87**, 071802 (2001).
[20] G. Burdman, E. Golowich, J. Hewett, and S. Pakvasa, Phys. Rev. D **66**, 014009 (2002).
[21] F. del Aguila and J. Santiago, Phys. Lett. B **493**, 175 (2000).