A Simple Flavor Protection for RS

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Abstract

We present a simple variation of warped flavor models where the hierarchies of fermion masses and mixings are still explained but dangerous flavor violating effects in the Kaon sector are greatly reduced. The key new ingredients are two horizontal U(1) symmetries. These symmetries forbid flavor violation in the down quark sector (with the exception of small IR brane localized kinetic mixing terms for the left-handed quarks) while allowing for flavor violation in the up quark sector. The leading flavor constraints come from $D^0 - \bar{D}^0$ mixing, and are safely satisfied for the KK mass scale of order 3 TeV. Our analysis of the flavor constraints also includes the constraints due to the (usually ignored) localized kinetic mixing terms. We also comment on the effects of the additional U(1) gauge bosons.
The Randall-Sundrum (RS) model [1] provides an interesting framework for new physics beyond the Standard Model (SM). Not only does it address the electroweak hierarchy problem because the Higgs sector localized on the IR brane has effectively a TeV cut-off scale, but it also offers an explanation of the SM fermion mass hierarchies via localized profiles of fermions in the extra dimension [2, 3]. As SM fermions get their masses from Yukawa interactions with the Higgs on the IR brane, one can localize the light fermions close to the UV brane to make their effective Yukawa couplings hierarchically small. The simplest “anarchic approach” assumes that the Yukawa couplings in the 5D theory are all $\mathcal{O}(1)$ in natural units. Then the localization of the SM fermions leads to the effective 4D mass matrices which naturally incorporate hierarchies both in the mass eigenvalues and the CKM mixing angles [4].

Theories of flavor quite generally predict new sources of flavor violation. In the RS model, flavor changing neutral currents (FCNC’s) arise already at the tree level, due to flavor violating couplings of the Kaluza-Klein (KK) modes to the quark mass eigenstates. These same KK modes play the important role of stabilizing the electroweak scale, and naturalness implies that the lightest KK mode should not be heavier than $\sim 3$ TeV. At first sight this spells disaster for the model, since tight experimental bounds on FCNCs would normally apply a much higher KK scale, $M_{KK} > 10^3 - 10^4$ TeV. However, the RS model has a built-in flavor protection usually referred to as RS-GIM [3, 5]: the localization mechanism responsible for the mass hierarchies also suppresses FCNC’s. RS-GIM is powerful enough to suppress below the experimental sensitivity almost all effective $\Delta F = 2$ four-fermion operators generated by the exchange of KK modes. The exception is the left-right (LR) operator $O_4 = (\bar{s}_L d_R)(\bar{s}_R d_L)$ whose imaginary part is tightly bounded by measurements of CP violation in the kaon sector. It turns out that the coefficient of this operator comes out roughly a factor of 100 too large for a $\sim 3$ TeV KK gluon [6], if one assumes no tuning among the input parameters. Thus, some form of flavor symmetry is required if we insist that the RS model addresses both the electroweak naturalness problem and the flavor hierarchies.

One possible approach is to require a full-fledged 5D GIM mechanism, where all tree-level FCNC’s are completely absent. This is indeed possible [7] using $U(3)^3$ flavor symmetries in the bulk, and assuming that the only source of flavor violation are fermion kinetic terms localized on the UV brane. The price is, however, that the hierarchy of fermion masses and mixings is left completely unexplained. One would prefer a solution that incorporates just enough flavor symmetries to eliminate the dangerous sources of flavor violation, while still allowing for the generation of flavor hierarchies. An attempt in this direction is the “5D MFV” proposal of [8], which postulates that only two spurions are responsible for breaking the flavor symmetry, generating both the bulk mass matrices and the brane Yukawa matrices. This assumption does not suppress FCNCs by itself and an additional alignment of bulk masses and brane Yukawa terms has to be imposed [8, 9]. Recently, an economical model based on a $U(3)$ symmetry acting in the bulk and broken on the IR brane was proposed in [10]. This symmetry, while allowing for down quark mixing via Yukawa couplings on the IR brane, enforces a degeneracy of the bulk masses in the down right quark sector which ensures universal couplings to vector KK modes. Flavor symmetries have also proved successful in warped models of the lepton sector, where they can explain the observed neutrino mixing
patterns and charged lepton flavor hierarchies without excessive lepton flavor violation [11].

In this paper we propose a set of symmetries for the quark sector that strongly protects the down quarks from flavor violations both in the right- and left-handed down sector, even after taking into account the effect of brane localized kinetic mixing terms. The idea is to use a horizontal $U(1)$ symmetry acting in the down quark sector. This symmetry acts in the bulk and on the IR brane, while it is broken on the UV brane. It aligns the bulk masses with the brane Yukawa matrices in the down sector: both are required to be diagonal in the basis in which the $U(1)$ charges are diagonal. This idea is similar to the SUSY alignment models [12], with the aim of moving all (or most) of flavor violation into the up sector. There is an immediate problem with this requirement: since the left-handed doublet quarks are charged under this $U(1)$, it seems we would need to assign charges in the right-handed up sector too to allow for the up-type Yukawa couplings. This would indeed be the case in the simplest RS-type model where only one bulk multiplet per generation hosts a quark doublet, and our $U(1)$ would forbid all flavor mixing and force the CKM matrix to be the unit matrix. However, one can construct well motivated models in which each SM quark doublet is embedded in several bulk multiplets. This is e.g. the case in the model of ref. [13] where the extended structure is introduced in order to ensure the custodial protection of the $Zb_Lb_L$ vertex [14]. In such extended models we are able to implement a horizontal $U(1)$ with desired properties. For reasons that will become clear in a moment we will need in fact two separate $U(1)$ symmetries, one of which is broken on the UV brane, and the other on the IR brane.

The model we propose is inspired by ref. [13], though here we deal with a smaller gauge symmetry group in the bulk and therefore we can utilize a more economical set of bulk fields. We consider 5D gauge theory in a slice of AdS$_5$. We parametrize the space-time by the conformal coordinates

$$ds^2 = \left( \frac{R}{z} \right)^2 (dx_\mu dx_\nu \eta^{\mu\nu} - dz^2),$$

where the AdS curvature is $R$, and the coordinate $z$ of the extra dimension runs between $R < z < R'$, $z = R$ corresponding to the UV (Planck) brane and $z = R'$ to the IR (TeV) brane. $R'/R \sim 10^{16}$ is the large number that sets the hierarchy between the Planck and the TeV scale. Our gauge group in the bulk is $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_X$ broken by boundary conditions on the UV brane down to $SU(3)_c \times SU(2)_L \times U(1)_Y$ [15]. The Higgs field $\Phi$ transforming as $(1, 2, 2)_0$ is localized on the IR brane. Each generation of

| $U(1)_q$ | $\Psi_u$ | $\Psi_e$ | $\Psi_d$ | $\Psi_{\ell}$ |
|-----------|-----------|-----------|-----------|-----------|
| $(+, -)$  | $u_1$     | $e_1$     | $d_1$     | $\ell_1$  |
| $(-, +)$  | $u_2$     | $e_2$     | $d_2$     | $\ell_2$  |

Table 1: Charge assignments under the horizontal $U(1)$ symmetries.
SM quarks is embedded in four 5D quark multiplets. The doublets are embedded in two bi-fundamental (under $SU(2)_L \times SU(2)_R$) quarks $\Psi_{q_u,q_d}$. This realization incorporates the custodial protection of the $Z\bar{b}_L b_L$ vertex [14] if the couplings of the two $SU(2)$’s are equal. The SM singlet quarks are embedded into $SU(2)_L \times SU(2)_R$ singlets $\Psi_{u,d}$. The representation under the full gauge group and the boundary conditions are chosen as

\[
\begin{align*}
(3,2,2)_{2/3} : \Psi_u = \begin{pmatrix} q_u[\pm,+] & \tilde{q}_u[-,+] \end{pmatrix} \\
(3,2,2)_{-1/3} : \Psi_d = \begin{pmatrix} \tilde{q}_d[-,+] & q_d[\pm,+] \end{pmatrix} \\
(3,1,1)_{2/3} : \Psi_u = \begin{pmatrix} u^c[-,-] \end{pmatrix} \\
(3,1,1)_{-1/3} : \Psi_d = \begin{pmatrix} d^c[-,-] \end{pmatrix}
\end{align*}
\]

In the square bracket we indicated the boundary conditions on the UV and IR branes, where $[+/-]$ denotes the right/left chirality of the bulk fermion vanishing on the brane. $[\pm]$ stands for mixed boundary conditions for the $SU(2)_L$ doublets:

\[
\theta q_{u,L}(0) - q_{d,L}(0) = 0 \quad q_{u,R}(0) + \theta^\dagger q_{d,R}(0) = 0 \tag{3}
\]

where $\theta$ is a $3 \times 3$ matrix. These boundary conditions could be obtained by coupling the combination $\theta q_{u,L} - q_{d,L}$ to three UV localized right-handed doublets $\psi_R$ via the mass term $\Lambda_0 \overline{\psi}_R(\theta q_{u,L} - q_{d,L})$ and then taking the limit of $\Lambda_0 \to \infty$. The effect of these boundary conditions is that the two bulk fields $q_{u,d}$ host only one left-handed zero mode $q_{L}(x)$. There are also right-handed zero modes $u_R(x), d_R(x)$ hosted by $u^c, d^c$. After electroweak symmetry breaking these zero modes obtain masses from the IR brane localized Yukawa interactions

\[
\mathcal{L}_{\text{yuk}} = -(R^4/R^3)\delta(z - R') \left\{ \text{Tr} \left[ \overline{\Psi}_{q_u,L} \Phi \right] \check{Y}_u \Psi_{u,R} + \text{Tr} \left[ \overline{\Psi}_{q_d,L} \Phi \right] \check{Y}_d \Psi_{d,R} \right\} + \text{h.c.} \tag{4}
\]

In the absence of flavor symmetries the Yukawa matrices are expected to be anarchical. We keep this assumption for the up-type Yukawa matrix, but we impose non-trivial structure on the down-type Yukawa matrix. We achieve this by introducing a $U(1)_d$ symmetry acting on the $X = -1/3$ quarks with generation dependent charges $(d_1, d_2, d_3)$, the same for $\Psi_{q_u}$ and for $\Psi_d$. There are no constraints on the values of the charges other than $d_i \neq d_j$ for $i \neq j$. The symmetry is valid in the bulk and on the IR brane but must be broken on the UV brane to allow for non-zero $\theta$. This symmetry will ensure that $c_{q_u}, c_{q_d}$ and the down Yukawa coupling $\check{Y}_d$ are simultaneously diagonal, and also forbids off-diagonal kinetic terms on the IR brane involving $\psi_{q_u}$ and $\psi_d$. The $U(1)_d$ can be global or local. In the latter case there is a corresponding gauge boson with ($+\$ boundary conditions. We will later comment on the phenomenology of such a gauge boson.

In this model there is another source of flavor violation in the down sector if the $\theta$ matrix that sets the UV boundary conditions is non-diagonal. To protect us from that we need to introduce another $U(1)_q$ symmetry that acts on the UV brane and in the bulk, but it has to be broken on the IR brane in order to allow non-diagonal up-type Yukawa couplings. Under this symmetry, $\Psi_{q_u}$ and $\Psi_{q_d}$ should have equal charges in each generation, but the charges of different generations should be different, $q_i \neq q_j$ for $i \neq j$. We can also charge the rest of the multiplets, which will have the effect of forbidding all off-diagonal brane kinetic terms on the UV brane. This symmetry will then ensure that $c_{q_u}, c_{q_d}$ and $\theta$ are all simultaneously
diagonal. Again, the symmetry can be global or local, in the latter case the $U(1)_q$ gauge boson has $(+−)$ boundary conditions.

The two $U(1)$ symmetries together ensure that $c_{qu}, c_{qd}, c_u, c_d, \bar{Y}_d$ and $\theta$ are all diagonal while allowing for off-diagonal components in $\bar{Y}_u$. Moreover, off-diagonal kinetic terms are not allowed on the UV brane (if we also charge $\psi_{u,d}$ under $U(1)_q$). On the IR brane, the symmetries still allow for two off-diagonal brane kinetic terms:

$$L = \frac{R^4}{R''^3} i\delta \left( \overline{\Psi}_{qu} \hat{K}_{qu} \phi \Psi_{qu} + \overline{\Psi}_{u} \hat{K}_{u} \phi \Psi_{u} \right) \delta(z - R) \quad (5)$$

We cannot decree that these terms vanish because they are generated by loop effects, but it is natural to assume that the coefficients are loop suppressed. It is important to note, that these kinetic mixing terms are present in every RS flavor model, but they are ignored because they are usually sub-leading. However, if the leading sources are suppressed this might turn out to be the dominant contribution. In particular in models with partial flavor protection [8, 10] these operators will be the sources of the leading flavor constraints. In the following we assume $\hat{K} \sim 1$ and the NDA value of the dimensionless coefficient

$$\delta \sim Y_*^2 \frac{\Lambda R''}{16\pi^2}. \quad (6)$$

where $Y_*$ is the typical amplitude of the entries in $\bar{Y}_{u,d}$, and $\Lambda$ is the cutoff scale for the Yukawa interaction, usually assumed to be around the second KK mode. The off-diagonal terms in $\hat{K}_{qu,u}$ lead to the kinetic mixing among the up-type quarks. This is a subleading effect compared to the mass mixing generated by the Yukawa terms so that it can be safely neglected. However, $\hat{K}_{qu}$ leads also to the kinetic mixing among the left-handed down quarks. This is a source of flavor violation in the down quark sector which we need to take into account. The improvement with respect to the standard RS flavor scenario is that the brane kinetic terms affect mostly the left-handed down quarks (so that there are no large LR currents) and are loop suppressed.

We are ready to analyze the flavor structure of our model. Let us write down the mass matrix for the SM quark fields $u_{L,R}, d_{L,R}$ which are zero modes of the bulk fermions. The localization of zero modes depends on the fermion bulk mass customarily written in terms of dimensionless $c$-parameters:

$$\frac{R^4}{z^5} \sum_{x=c_{qu}, c_{qd}, c_u, c_d} \overline{\Psi}_x c_x \Psi_x. \quad (7)$$

where each $c$ is a $3 \times 3$ matrix in the flavor space, $c_x = \text{diag}(c^1_x, c^2_x, c^3_x)$. The zero modes are
embedded into the 5D fields as

\[
q_{u,L}(x, y) \rightarrow R'^{-1/2} \left( \frac{z}{R} \right)^2 \left( \frac{z}{R'} \right)^{-c_{qu}} f_{qu} C_\theta q_L(x)
\]

\[
q_{d,L}(x, y) \rightarrow R'^{-1/2} \left( \frac{z}{R} \right)^2 \left( \frac{z}{R'} \right)^{-c_{qd}} f_{qd} S_\theta q_L(x)
\]

\[
u_R^u(x, y) \rightarrow R'^{-1/2} \left( \frac{z}{R} \right)^2 \left( \frac{z}{R'} \right)^{c_u} f_{-u} u_R(x)
\]

\[
d_R^d(x, y) \rightarrow R'^{-1/2} \left( \frac{z}{R} \right)^2 \left( \frac{z}{R'} \right)^{c_d} f_{-d} d_R(x)
\]

(8)

where \( f_x = \text{diag}(f(c_{x1}), f(c_{x2}), f(c_{x3})) \) and \( f(c) \) is defined as

\[
f(c) = \frac{\sqrt{1 - 2c}}{\left[1 - (\frac{R'}{R})^{2c-1}\right]^{1/2}}.
\]

(9)

Furthermore, we traded \( \theta \) for the diagonal matrices \( C_\theta, S_\theta \) defined as

\[
C_\theta = \frac{1}{\sqrt{1 + |\hat{\theta}|^2}} \quad S_\theta = \frac{\hat{\theta}}{\sqrt{1 + |\hat{\theta}|^2}} \quad \hat{\theta} = \theta \frac{f_{qu}(R'/R)^{c_{qu}-c_{qd}}}{f_{qd}}
\]

(10)

which decide how the doublet zero modes are distributed between the two bulk multiplets. In this basis the kinetic terms for the zero-mode quarks are not diagonal due to the brane kinetic terms in eq. (5). We can safely neglect the kinetic mixing of the up quarks since the mixing via the Yukawa couplings is much larger, but we should take into account the kinetic mixing of the down quarks which is the leading effect. The kinetic terms for the left-handed quark are given by,

\[
K = 1 + \delta C_\theta f_{qu} \tilde{K}_q f_{qu} C_\theta.
\]

(11)

To diagonalize them, we need to perform the Hermitian transformation \( d_L(x) \rightarrow H d_L(x) \) where \( HKH = 1 \). After that, we can plug the zero mode profiles into the brane Yukawas and replacing the Higgs with its vev to obtain the SM mass matrices,

\[
\mathcal{M}^{SM}_u = \frac{v}{\sqrt{2}} C_\theta f_{qu} \tilde{Y}_u f_{-u}
\]

\[
\mathcal{M}^{SM}_d = \frac{v}{\sqrt{2}} H S_\theta f_{qd} \tilde{Y}_d f_{-d}
\]

(12)

To proceed, we need to find the bi-unitary rotations that diagonalize these mass matrices. We start with the up quark mass matrix. The natural expectation is that \( \tilde{Y}_u \) is anarchical, and below we work under this assumption. On the other hand, as there are no non-abelian flavor symmetries to enforce the equality of the c-parameters, \( f_{qu,qd,-u,-d} \) are expected to
be hierarchical. Below we assume \( f_{u1} \ll f_{u2} \ll f_{u3} \) and \( f_{u1} \ll f_{u2} \ll f_{u3} \). Then the up quark matrix can be diagonalized as \( \mathcal{M}_u^{SM} = L^u(m_u)_{\text{diag}} R_u^\dagger \), where the elements of the rotation matrices are approximately given by

\[
|L_{ij}| \sim \frac{f_{u1}}{f_{u3}}, \quad |R_{ij}| \sim \frac{f_{u1}}{f_{u3}}, \quad i \leq j.
\]

Since the dominant source of flavor mixing comes from the up sector, the CKM matrix is approximately given by the hermitian conjugate of the rotation matrix \( L_u \). Thus the hierarchy in the CKM matrix elements is set by the \( c_{q_u} \) parameters and we need to choose them such that

\[
\frac{f_{u1}}{f_{u2}} \sim \lambda^2, \quad \frac{f_{u2}}{f_{u3}} \sim \lambda^3,
\]

where \( \lambda \sim \sin \theta_C \sim 0.2 \). The remaining \( c \) parameters have to be adjusted to reproduce the quark masses. In this model we have actually more \( c \) parameters than quark masses to fit so several options are possible. For the purpose of illustration, we propose here one particular pattern. We assume that all \( c \) parameters for the third generation are such that the zero modes are localized in the IR, that is \( c_{q_u,d}^3 < 1/2, c_{q_u,d}^3 - 1/2 \), whereas the first two generations are localized in UV: \( c_{q_u,d}^{1,2} > 1/2, c_{q_u,d}^{1,2} < -1/2 \). At this point, important observables depend on the relation between \( c_{q_u} \) and \( c_{q_d} \). We define three dimensionless numbers

\[
\alpha_i = (R_i^d/R_i^u)_{q_u} - c_{q_d}
\]

As we will see in the moment, in this scenario the \( m_b/m_t \) ratio is set by \( |\alpha_3| \). The small ratio can be obtained by choosing \( c_{q_u}^3 < c_{q_d}^3 \) so as to make \( |\alpha_3| \ll 1 \), rather then localizing the right bottom in UV as in the standard approach to RS flavor. This implies that \( \hat{\theta}_3 \ll 1 \), whereas \( \hat{\theta}_1,2^2 \sim 1 \), as the ratio of \( f \)’s and the exponent cancel in the later case. For the moment we leave the options open for \( \alpha_{1,2} \sim f_{q_u}^{1,2}/f_{q_u}^{1,2} \): they can be smaller or larger than 1.

Next, we discuss the down quark mass matrix. Although the down-type Yukawa couplings are diagonal, the mass matrix is slightly tipped by the Hermitian transformation \( H \). The off-diagonal terms in the kinetic matrix are suppressed by both a loop factor and by RS-GIM and, in consequence, the Hermitian rotation matrix is close to the unit matrix \( (H)_{ij} \sim \delta_{ij} - \frac{\phi}{2} f_{q_u} f_{q_u} \). The unitary rotations \( L_d^\dagger \) and \( R_d^\dagger \) that diagonalize down mass matrix have smaller off-diagonal terms than the corresponding up quark rotations,

\[
|L_{ij}^d| \sim \frac{\delta}{2} f_{q_u} f_{q_u}^i, \quad |R_{ij}^d| \sim \frac{\delta}{m_d^2} f_{q_u} f_{q_u}^i, \quad i < j.
\]

The eigenvalues of the quark mass matrices are of the order of

\[
(m_u)_{\text{diag}} \sim \frac{vY_u}{\sqrt{2}} f_{q_a} f_{-u}, \quad (m_d)_{\text{diag}} \sim \frac{vY_u}{\sqrt{2}} f_{q_u} f_{-d}.
\]

Thus, the bottom to top quark mass ratio is \( m_b/m_t \sim \alpha_3 f_{-d^3} / f_{-u^3} \) which is of order \( \alpha_3 \) as long as \( f_{-d^3} \sim f_{-u^3} \sim 1 \). In order to reproduce the up quark masses, the ratios of \( f_{-u} \)’s
should be adjusted as
\[ \frac{f_{-u^1}}{f_{-u^3}} \sim \frac{m_u}{m_t} \lambda^{-3}, \quad \frac{f_{-u^2}}{f_{-u^3}} \sim \frac{m_c}{m_t} \lambda^{-2}. \tag{18} \]
This fixes the magnitude of the rotation angles in the right-handed up quark sector. As for the remaining down quark masses, we need
\[ \alpha_1 \frac{f_{-d^1}}{f_{-u^3}} \sim \frac{m_d}{m_t} \lambda^{-3}, \quad \alpha_2 \frac{f_{-d^2}}{f_{-u^3}} \sim \frac{m_s}{m_t} \lambda^{-2}, \tag{19} \]
we can adjust \( \alpha_{1,2} \) or \( c_d^{1,2} \) or both to match the experimental values.

We move to discussing the FCNCs in our model. Typically, the strongest constraints come from \( \Delta F = 2 \) processes mediated by tree-level exchange of the KK gluons, as their coupling to matter fields is usually stronger than that of other vector KK modes. Thus the first task is to determine the off-diagonal couplings of the KK gluon to the zero-mode quarks. In the original flavor basis the coupling depends on the parameter \( c \) of the zero mode, and is approximately given by \( [5, 6] \)
\[ g_x \approx g_{ss} \left( -\frac{1}{\log R'/R} + f_x^2 \gamma(c_x) + \delta f_x \tilde{K}_x f_x \right). \tag{20} \]
where \( g_{ss} \) is the bulk \( SU(3) \) gauge couplings \( (g_{ss} \sim 6 \) in the absence of brane kinetic terms for the \( SU(3) \) gauge fields), and \( \gamma(c) \) is of order one but it varies by order 1 values with \( c \) varied in the range of interest. The last term is due to the IR brane kinetic terms for the fermions. Non-diagonal couplings arise if \( \tilde{K}_x \) is non-diagonal, or they may be generated by the unitary transformation from the flavor basis to the mass-eigenstate basis if \( f_x \) is non-universal. In our model both of these sources are present.

The leading off-diagonal couplings to the left- and right-handed up quarks can be read from
\[ (g_{uL}) \sim g_{ss} L_u^\dagger f_{q_i}^2 (1 + \alpha^2) L_u \quad (g_{uR}) \sim g_{ss} R_u^\dagger f_{-u^3} R_u. \tag{21} \]
Performing the rotation, we find that the off-diagonal couplings of the KK gluon to right handed quarks are given by \( c_u \), \( (g_{R})_{ij} \sim f_{-u^3} f_{-u^j} \), which in turn depend on the observable CKM mixing angles and mass ratios and on \( f_{-u^3} \sim 1 \). The RS-GIM mechanism is at work in the sense that the couplings to the light quarks are suppressed by the hierarchically small \( f_{-u} \). In particular, for the up and the charm quarks,
\[ (g_{uR})_{12} \sim g_{ss} \frac{m_u m_c}{m_t^2} \frac{f_{-u^3}}{\lambda^5} \sim g_{ss} \frac{2m_u m_c}{Y_{u^2} v^2} \frac{1}{\lambda^5 f_{q_d}^2}, \tag{22} \]
which is suppressed by \( m_u m_c/m_t^2 \lambda^5 \sim 5 \cdot 10^{-5} \). For the couplings of the left-handed quarks we have some more freedom. If all \( |\alpha_i| \ll 1 \) then \( (g_{uL})_{ij} \sim g_{ss} f_{q_i} f_{q_d} \) and the couplings can be expressed in terms of the CKM hierarchy. For \( |\alpha_{1,2}| > 1 \) the couplings may be enhanced. For the up and the charm quarks we find
\[ (g_{uL})_{12} \sim g_{ss} f_{q_1} f_{q_2} \left( 1 + \alpha_2^2 + \alpha_3^2 + \frac{J_{q_1}^2}{J_{q_2}^2} \alpha_1^2 \right) \sim g_{ss} \lambda^5 f_{q_3}^2 (1 + \alpha_2^2) \tag{23} \]
which is RS-GIM suppressed by $\lambda^5 \sim 6 \cdot 10^{-4}$. In the second step we used our assumptions $|\alpha_3| \ll 1$, $f_{q_{1}^4} \ll f_{q_{2}^4}$. Writing $1 + \alpha^2$ we really mean the larger of the two numbers, since we are estimating a sum of complex contributions with random phases.

For the left-handed down quark, the off-diagonal KK gluon couplings appear due to the brane-kinetic terms even before diagonalizing the kinetic terms and the mass matrix. After performing the Hermitian rotation, the couplings become

$$(g_{d_L}) \sim g_{ss} H \left( -\frac{1}{\log(R'/R)} + f_{q_u}^2 (1 + \alpha^2) + \delta f_{q_u} \tilde{K}_{q} f_{q_u} \right) H$$

(24)

One can see that the matrix $H$, designed to diagonalize the kinetic terms, does not simultaneously diagonalize the KK gluon couplings. Instead, we are left with the off-diagonal terms $(g_{d_L}^i)_{ij} \sim g_{ss} \delta f_{q_u} f_{q_u}$. The subsequent unitary rotation $L_d$ does not generate any larger terms. For the down and strange quarks we find

$$(g_{d_L})_{12} \sim g_{ss} \delta \lambda^5 f_{q_u}^2.$$  

(25)

Besides the RS-GIM suppression, this coupling enjoys the loop suppression factor $\delta$. For the right-handed down quarks, the unitary rotation $R_d$ is the only source of flavor off-diagonal couplings who for this reason become even more suppressed: $(g_{d_R})_{ij} \sim g_{ss} \delta (m_{d}/m_{d}) f_{q_u} f_{q_u} f^2_{-d}$, $i < j$. For the down and strange quarks

$$(g_{d_R})_{12} \sim g_{ss} \delta \frac{m_{d}}{m_{s}} \lambda^5 f_{q_u}^2 f^2_{-d} \sim g_{ss} \delta \frac{2 m_{d} m_{s}}{Y_s v^2} \frac{\lambda}{\alpha_2^2}.$$  

(26)

Having obtained the flavor violating couplings, we can estimate the magnitude of the flavor-violating four-fermion operators generated by a tree-level exchange of the KK gluons. After applying the Fierz identities the effective Hamiltonian is of the form [6]

$$\mathcal{H} = \frac{1}{M_C} \left[ \frac{1}{6}(g_{q_L})_{ij}(g_{q_L})_{kl}(q_{L}^{i\alpha} \gamma_\mu q_{L}^{j\mu}) (q_{L}^{k\beta} \gamma_\mu q_{L}^{l\beta}) + \frac{1}{6}(g_{q_R})_{ij}(g_{q_R})_{kl}(q_{R}^{i\alpha} \gamma_\mu q_{R}^{j\mu}) (q_{R}^{k\beta} \gamma_\mu q_{R}^{l\beta}) 
- (g_{q_R})_{ij}(g_{q_L})_{kl} \left( (q_{R}^{i\alpha} q_{R}^{k\beta}) (q_{L}^{j\mu} q_{R}^{l\beta}) - \frac{1}{3} (q_{R}^{i\alpha} q_{R}^{k\beta}) (q_{L}^{j\mu} q_{L}^{l\beta}) \right) \right]$$

(27)

where $\alpha, \beta$ are color indices. This parameterization of flavor violating new physics contributions in terms Wilson coefficients $C(\mu)$:

$$\mathcal{H} = C^1(\mu)(q_{L}^{i\alpha} \gamma_\mu q_{L}^{j\mu}) (q_{L}^{k\beta} \gamma_\mu q_{L}^{l\beta}) + C^1(\mu)(q_{R}^{i\alpha} \gamma_\mu q_{R}^{j\mu}) (q_{R}^{k\beta} \gamma_\mu q_{R}^{l\beta})
+ C^4(\mu)(q_{R}^{i\alpha} q_{R}^{k\beta}) (q_{L}^{j\mu} q_{L}^{l\beta}) + C^5(\mu)(q_{R}^{i\alpha} q_{R}^{k\beta}) (q_{L}^{j\mu} q_{R}^{l\beta})$$

is subject to constraints from low-energy experiments [16]. In our model the strongest constraints come from $\Delta F = 2$ transitions between the up and the charm quarks in $D^0 - \bar{D}^0$ mixing. The corresponding Wilson coefficients are given by

$$C^1_D(M_G) \sim \frac{1}{M_C^2} \frac{g_{ss}^2}{6} \lambda^{10} f_{q_u}^4 (1 + \alpha_2^2)^2 \quad \tilde{C}^1_D(M_G) \sim \frac{1}{M_C^2} \frac{g_{ss}^2}{2} \frac{m_u^2 m_c^2}{3v^4} \frac{1}{\lambda^{10} f_{q_u}^4}$$

(28)
Table 2: Flavor violation in the D- and K-meson sector. In the second column we give the suppression scale of the Wilson coefficients in our $U(1)$ protected model (for $\alpha_2 \sim 1$). The numerical value of the suppression scale is evaluated for two representative choices $f_{q_d}^2$, and for $g_{ss} = 6, \lambda_2 = 1, \delta = Y_s^2/16\pi^2 (\Lambda R^2) \approx Y_s^2/4\pi^2$. In the last column we give the model independent lower bound on the suppression scale from ref. [16].

$$
C_D^i(M_G) = -3C_D^i(M_G) \sim \frac{1}{M_G^2 Y_s^2} \left(1 + \alpha_2^2\right) (1 + \alpha_2^2)
$$

We also give the Wilson coefficients in the kaon sector

$$
C_K^1(M_G) = \frac{1}{M_G^2} \frac{g_{ss}^2}{6} \lambda^2 f_{q_d}^4 \quad C_K^1(M_G) \sim \frac{1}{M_G^2} \frac{g_{ss}^2}{6} \lambda^2 f_{q_d}^4
$$

$$
C_K^3(M_G) = -3C_K^3(M_G) \sim \frac{1}{M_G^2 Y_s^2} \left(1 + \alpha_2^2\right) (1 + \alpha_2^2)
$$

We now have everything we need to discuss flavor constraints on our model. Inserting the experimentally measured values of the quark masses and mixing angles at the TeV scale (see Table 1 in [6]), and making some assumptions concerning $g_{ss}, Y_s$ and $c_{q_d}^2$, we can find how much are the Wilson coefficients suppressed with respect to $1/M_G^2$. Comparing it with the experimental lower bound we obtain a constraint on $M_G$, that is on the KK scale. Sample results are given in Table 2 where we took $g_{ss} = 6, Y_s = 1$. In the kaon sector, the LR operator (who is the most troublesome in the standard RS flavor scenario with anarchic flavor in the down sector) comes out very suppressed in our model and does not pose any problems. As for the LL operator, it satisfies the experimental bounds even in the standard RS with a 3 TeV KK gluon. In our case, the LL operator is down by the additional loop

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In the standard RS flavor scenario one usually takes $Y_s \sim 3$, close to the perturbativity bounds, in order to maximally suppress the dangerous FCNC’s. In our case the flavor bounds can be satisfied with Yukawa couplings in a safely perturbative region. The loop induced kinetic mixings are then also naturally suppressed since in this case $\delta = Y_s^2/16\pi^2 (\Lambda R^2) \approx Y_s^2/4\pi^2 \ll 1$. 

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factor $\delta^2$ and thus is always well below the experimental limits. That contribution could be sizable only if $\delta \sim 1$, that is for large IR brane kinetic terms.

Quite unlike the standard RS flavor scenario, the strongest bounds come from the D-meson sector. In Table \[2\] we present our results for two different choices for the localization of the 3rd generation doublet quark. In both cases we take $c_{q_u}^3 < 1/2$ corresponding to the IR localization, but in the first case ($f_{q_u}^3 = 0.3$) the profile is fairly flat, while in the second case ($f_{q_u}^3 = 0.3$) it is sharply localized in IR. The strongest bound always comes from the LR current, that is from the coefficient $C_{\text{LL}}^1$, but as long as $Y_{s} > 1$ a 3 TeV KK gluon is allowed (a larger Yukawa coupling leads to relaxing this bound). In the case with $f_{q_u}^3 = 1$ another constraint comes from the LL operator yielding a 2 TeV bound on $M_G$. In any case, the bounds are weaker than the ones from electroweak precision tests that push the KK scale to around 3 TeV anyways [15]. The reason is that the flavor constraints in the D meson sector are less severe than analogous constraints in the kaon sector since $D^0 - \bar{D}^0$ mixing is more challenging concerning both the theoretical prediction in the SM and the actual measurement (see e.g. [17]). Thus, imposing flavor symmetries in the down sector and leaving an anarchic flavor violation in the up sector leads to very mild flavor bounds on the scale of new physics.

Some more constraints on our model may be obtained by investigating $\Delta F = 1$ processes. Those generated at the loop level are beyond the scope of this paper, however there are also tree-level $\Delta F = 1$ FCNC transitions mediated by the electrically neutral KK modes of $SU(2)_L \times SU(2)_R$ gauge bosons (the KK Z). The off-diagonal couplings of KK Z to the left-handed quarks are analogous to those of the KK gluon: $(g_Z^2)_{ij} \sim (\pm 1/2)g_{ws}f_{q_i}^2f_{q_j}^2$, where $g_{ws}$ is the $SU(2)$ bulk gauge coupling ($g_{ws} \sim 4$ in the absence of brane kinetic terms). Electroweak symmetry breaking mixing the KK Z with the SM Z boson via the IR brane localized Higgs, with a mixing mass of order $g_{ws}m_Zv/2$ for both $SU(2)_L$ and $SU(2)_R$ gauge bosons. This mixing induces the effective coupling of the left-handed up quarks to the Z boson, for example $(g_Z^2)_{ct} \sim (g_{ws}^2/4)\lambda^2f_{q_i}^2m_Zv/M_{Z'}^2$, where $M_{Z'}$ is the mass of the lightest KK Z. Note that for the down-type quarks $(g_L^2)_{ba}$ is much more suppressed because the $L$ and $R$ KK exchange approximately cancels in that case. The right-handed quarks, on the other hand, live in the $SU(2)_L \times SU(2)_R$ singlets so that their effective couplings to the Z boson are also suppressed. To describe the effective $t_LZ_{CL}$ coupling, the low energy theory contains the dimension six operator $(\tilde{H} = i\sigma_2H^*)$

$$O_{LL}^u = i \left[ \overline{Q}_3\tilde{H} \right] \left[ (\overline{\psi}\tilde{H}^\dagger)Q_2 \right] - i \left[ \overline{Q}_3(\overline{\psi}\tilde{H}) \right] \left[ \tilde{H}^\dagger Q_2 \right] + \text{h.c.}$$

with the coefficient

$$C_{LL}^u \sim \frac{1}{2M_{Z'}^2}g_{ws}^2\lambda^2f_{q_i}^2 \sim 0.04 \left( \frac{g_{ws}}{4} \right)^2 \left( \frac{f_{q_i}}{1} \right)^2 \left( \frac{3\text{ TeV}}{M_{Z'}^2} \right)^2 \text{ TeV}^{-2}$$

We took into account the fact the $L_\mu$ and $R_\mu$ have almost degenerate KK modes, and we assumed both the $L$ and $R$ coupling to be $g_{ws}$. The coefficients of the analogous operators involving an up quark are more suppressed by RS-GIM. Now, at the loop level the operator
in eq. (32) feeds in to $B \rightarrow X_s \gamma$ and $B \rightarrow X_s l^+ l^-$, which leads to the bound $|C_{LL}^S| < 0.07 \text{ TeV}^{-2}$ [21]. This is satisfied for a 3 TeV KK scale, as $f_{q_u}$ is not much larger than 1.

If $U(1)_d \times U(1)_q$ is promoted to a local symmetry (as expected from the AdS/CFT correspondence), there will be the corresponding KK gauge bosons which couple to the SM quarks. In case that the $U(1)$ gauge bosons are present, they will impose additional constraints on the parameters of the model. First of all, we have to worry about the limits from direct searches for new light neutral gauge bosons (aka Z'). The $U(1)_d$ KK modes do not pose a problem. As a consequence of the [−+] boundary condition, they are localized in the IR and their masses are at the KK scale which we assume to be at least 3 TeV. On the other hand, the lightest $U(1)_q$ gauge boson has an almost flat profile (except near the IR brane) and its mass is suppressed with respect to the KK scale (much as the W and Z masses in Higgsless models [18]): $M_q \sim (1/R') \sqrt{2/\log(R'/R)} \sim .1 M_G$ which is around 300 GeV when the lightest KK gluon is at 3 TeV. One way to evade the Tevatron limits is to assume that $U(1)_q$ does not couple to the SM leptons, another is to assume that the $U(1)_q$ gauge coupling is small enough. In fact, the smallness of the coupling is also required by the flavor bounds. Since the $U(1)_q$ couplings to left-handed quarks is necessarily non-universal, the rotation of the up-type quarks to the mass eigenstate basis generates flavor non-diagonal couplings

\[ (g_L^q)_{ij} \sim \frac{g_{q^*}}{\log^{1/2}(R'/R)} |q_i - q_j| \frac{f_{q_u}}{f_{q_d}} i < j \quad (34) \]

where $g_{q^*}$ is the $U(1)_q$ bulk coupling, and the coupling of the lightest KK mode is approximately $g_{q^*}/\log^{1/2}(R'/R)$. Unlike the KK gluon couplings, the flavor violating coupling to $u$ and $c$ is suppressed by only one power of the Cabibbo angle. The RS-GIM mechanism is not acting here because this gauge boson is not IR brane localized, and the SM fermions have family dependent couplings. This fact together with the unbearable lightness of the $U(1)_q$ gauge boson implies that $g_{q^*}$ coupling must be much smaller than the strong bulk coupling to pass the flavor bounds. More precisely, the tree-level exchange of the $U(1)_q$ gauge boson contributes to the Wilson coefficient $C_D$ as

\[ C_D^1 \sim \frac{g_{q^*}^2}{2\log(R'/R)M_q^2} \lambda^2 \sim \frac{g_{q^*}^2}{(12 \text{ TeV})^2} \left( \frac{3 \text{ TeV}}{M_G} \right)^2 \quad (35) \]

Note, that a Fierz factor of 3 in the denominator is missing because a U(1) rather than SU(3) gauge boson is exchanged. We can see that we need $g_{q^*} < 0.01$ to pass the experimental bounds listed in Table 2 with a 3 TeV KK scale. Thus the $U(1)_q$ symmetry is required to be global for all practical purposes.

On the other hand, the bounds on the $U(1)_d$ gauge coupling $g_{d_u}$ are less severe, even though the couplings of these KK modes do not have a full RS-GIM protection either. The reason is that the $U(1)_d$ gauge bosons would be localized in the IR just like the KK gluon,

\[ ^3 \text{We thank Uli Haisch for pointing this out to us.} \]

\[ ^4 \text{However, one may as well use two discrete symmetries } Z_N \times Z_M \text{ with charges and symmetry breaking patterns identical to our } U(1) \times U(1) \text{ symmetries. In that case we would not have to introduce } U(1) \text{ gauge bosons. We thank Yuval Grossman for this suggestion.} \]
but the charges are now generation dependent so that the rotation to the mass eigenstate basis leads to flavor violating couplings. The largest contribution corresponds to the first log suppressed term in (20), which could be ignored for the KK gluon since it is generation independent in that case. After the rotation of the left-handed up quarks, the leading non-diagonal couplings are
\[ \sim g_d^* \log \frac{R'}{R} L_{ij}^u |d_i \alpha_i^2 - d_j \alpha_j^2| \]
and the contribution to the Wilson coefficient \( C_D^1 \) is given by
\[ C_D^1 \sim \frac{g_d^*}{2M_G^2} \frac{2}{\log R' / R} \alpha_i^4 \sim \frac{g_d^*}{(700 \text{ TeV})^2} \alpha_i^4 \left( \frac{3 \text{ TeV}}{M_G} \right)^2. \] (36)

There could be an additional suppression by \( \alpha_i \) because only quark component in the 5D multiplet \( \Psi_{qd} \) is charged under \( U(1)_d \), and the left-handed up-type quarks could live dominantly in \( \Psi_{qu} \) if the \( \alpha_i \) are small. Without the additional \( \alpha_i \) suppressions one would need \( g_d < 1/2 \) (somewhat weak), while with \( \alpha_i \) suppression \( g_d \) may well be of the similar strength as the strong and electroweak bulk couplings.

Let us now discuss the experimental signatures. The model we proposed predicts new physics contributions to FCNC’s in the D-meson sector at the level close to the current experimental bound. Improvement in both theory and experiment could lead to pinpointing new physics in the D-meson mixing and thus provide the first hint toward our model. The model also predicts loop induced contributions to \( \Delta F = 1 \) processes in the down sector, for example to \( B \rightarrow X_s \gamma \). On the other hand, a discovery of new sources of flavor and CP violation in the K- or B-meson mixing would – if not rule out the model – force it into an awkward region of parameter space with large brane kinetic terms. At the LHC, the RS framework obviously predicts spectacular new phenomena in abundance: KK modes of the SM fields, new gauge bosons and exotic fermions required by the custodial symmetry, see [19] for some recent papers. Our flavor model does not add many new states to that lot, though we may hope for a \( \sim 3 \text{ TeV} \) \( U(1)_d \) gauge boson with non-universal couplings to the SM quarks. Some more hints from the LHC may be revealed by FCNC processes involving the top quark where the RS flavor models are known to significantly enhance the SM predictions [20]. One sensitive probe is the decay process \( t \rightarrow cZ \), which is observable at the LHC at the level of \( 10^{-5} \) branching ratio. However, in our model this decay proceeds mostly via left-handed quarks, and there is the tension between obtaining a large enough coefficient \( C_{LL}^u \) and the indirect constraint from \( B \rightarrow X_s \gamma \) [21]. This makes the prospects for a signal at the LHC challenging.

Finally, we comment on extending our idea to models of gauge-higgs unification, where the Higgs boson is a pseudo-Goldstone boson whose mass is protected by approximate global symmetries. For example, in the model of ref. [13] based on the \( SU(3) \times SO(5) \times U(1) \) gauge group in the bulk, the SM quark doublets also are embedded in two different bulk multiplets, so that our \( U(1) \) protection can be implemented. An important difference with respect to our RS model is that in gauge-higgs unification the kinetic mixing feeds the flavor violation into the left-handed down quark sector even in the absence of brane kinetic terms [6]. Instead, the kinetic mixing originates from the same source as the Yukawa couplings (which is the boundary mass terms), therefore it is not loop suppressed as in our RS model. Nevertheless, in the limit of unbroken electroweak symmetry there is no flavor violation in the right-handed
down sector which suppresses the most dangerous contributions to the LR four-fermion operators in the kaon sector. The gauge-higgs model thus displays a similar phenomenology to the model of ref. [10], with sizable flavor violation in both up and down quark sectors.

In summary, we have presented a warped flavor model where most of the flavor violation is moved into the up sector. The key elements are introducing two horizontal U(1) symmetries, and embedding the LH quark doublets into two separate bulk multiplets. As a result, all bulk mass parameters are aligned with the down-type Yukawa coupling $\bar{Y}_d$ and the extra mixing parameter $\theta$ (the latter is needed due to the double embedding of the LH quarks). The only source for flavor violation in the down sector are small IR brane localized kinetic mixing terms. As a consequence, the leading constraints arise from charm (D) physics (rather than K or B physics), but these are also all satisfied with a KK mass scale of 3 TeV. If our U(1)’s are gauged, then these new gauge bosons will introduce new sources of flavor violation, setting stringent bounds on the possible sizes of these new gauge couplings.

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