Physics Letters B 742 (2015) 178–182

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Unified flavor symmetry from warped dimensions

Mariana Frank [a,*], Cherif Hamzaoui [b], Nima Pourtolami [a], Manuel Toharia [a]

[a] Department of Physics, Concordia University, 7141 Sherbrooke St. West, Montreal, Quebec, H4B 1R6, Canada
[b] Groupe de Physique Théorique des Particules, Département des Sciences de la Terre et de L'Atmosphère, Université du Québec à Montréal, Case Postale 8888, Succ. Centre-Ville, Montréal, Québec, H3C 3P8, Canada

ARTICLE INFO

Article history:
Received 31 October 2014
Received in revised form 14 January 2015
Accepted 21 January 2015
Available online 23 January 2015
Editor: M. Cvetič

Keywords:
Warped extra dimensions
Neutrinos
Fermion masses

ABSTRACT

In a model of warped extra-dimensions with all matter fields in the bulk, we propose a scenario which explains all the masses and mixings of the SM fermions. In this scenario, the same flavor symmetric structure is imposed on all the fermions of the Standard Model (SM), including neutrinos. Due to the exponential sensitivity on bulk fermion masses, a small breaking of this symmetry can be greatly enhanced and produce seemingly un-symmetric hierarchical masses and small mixing angles among the charged fermion zero-modes (SM quarks and charged leptons), thus washing out visible effects of the symmetry. If the Dirac neutrinos are sufficiently localized towards the UV boundary, and the Higgs field leaking into the bulk, the neutrino mass hierarchy and flavor structure will still be largely dominated and reflect the fundamental flavor structure, whereas localization of the quark sector would reflect the effects of the flavor symmetry breaking sector. We explore these features in an example based on which a family permutation symmetry is imposed in both quark and lepton sectors.

1. Introduction

The original motivation for warped extra-dimensions, or Randall–Sundrum models (RS), was to address the hierarchy problem. In RS the fundamental scale of gravity is exponentially reduced from the Planck mass scale to a TeV size due to a Higgs sector localized near the boundary of the extra dimension [1]. If SM fermions are allowed to propagate in the extra dimension [2], and become localized towards either boundary, the scenario also addresses the flavor problem of the SM and suppresses generic flavor-violating higher-order operators present in the original RS setup. However, KK-mediated processes still generate dangerous contributions to electroweak and flavor observables (including dangerous deviations to the $Zb\bar{b}$ coupling) [3–5], pushing the KK scale to 5–10 TeV [6]. The usual mechanisms employed to lower the KK scale involve using a custodial gauge $SU(2)_R$ symmetry [7], which insures a small contribution to electroweak precision parameters, lowering the KK scale to around 3 TeV, and rendering this scenario visible at the LHC. Alternatively, introducing a dilatonic scalar such that the warping of the 5-th dimension is strongly modified near the infrared (IR) brane while behaving as an Anti-de-Sitter space near the ultraviolet region, (so-called soft-wall metrics) is another solution [8]. There, the hierarchy problem imposes constrains on the Higgs profile which are stronger than in the original model, while electroweak constraints are milder, allowing a KK scale as low as 1–3 TeV.

One realization of the warped space model is based on the so-called flavor anarchy [5], in which one assumes that no special structure governs the flavor of Yukawa couplings and bulk fermion masses, as natural $\mathcal{O}(1)$ values for these 5D parameters already generate viable masses and mixings. However the neutrino sector must behave differently, first due to the possibility of Majorana mass terms, and second because this setup generates large mass hierarchies and small mixing angles, at odds with neutrino observations. An interesting property of flavor anarchy warped scenarios was investigated in [9], for the case of a bulk Higgs wave function leaking into the extra dimension. There one would obtain small mixing angles and hierarchical masses for all charged fermions, and at the same time very small Dirac masses, with large mixing angles and negligible mass hierarchy for neutrinos. Thus the flavor anarchy paradigm could still work in these scenarios. The case of Majorana neutrinos in this setup was also addressed in [9] where small and almost degenerate masses can be generated for IR localized Majorana neutrinos, out of higher-order interactions in the Lagrangian. The problem is that the warping necessary to

* Corresponding author.
E-mail addresses: mariana.frank@concordia.ca (M. Frank), hamzaoui.cherif@uqam.ca (C. Hamzaoui), n_pour@live.concordia.ca (N. Pourtolami), mtoharia@physics.concordia.ca (M. Toharia).

http://dx.doi.org/10.1016/j.physletb.2015.01.025

obtain the correct neutrino masses would be too small, and not enough to solve the hierarchy problem. Adding UV-localized Majorana masses (via some $SU(2)$ triplet representation in the UV region) can also generate small and almost degenerate masses [10]. As our goal here is to treat neutrinos on equal footing as quarks and charged leptons, we prefer not to add neutrino-specific extensions and assume that Dirac neutrino masses give the dominant contribution to neutrino masses.

Here we present an alternative scenario where instead of adopting flavor anarchy, we propose that all fermions share the same flavor symmetry. We assume that the flavor violating effects in the 5D Lagrangian can be parametrized by a small coefficient whose size is controlled by a ratio of scales, $\frac{\langle\phi\rangle^n}{\Lambda^n}$, with $\langle\phi\rangle$ the vacuum expectation value (VEV) of some flavon field, and Λ some cut-off mass scale, or the KK mass of some other flavon fields. This small breaking of the flavor symmetry is enough to reproduce correctly the flavor structure of the SM in both the quark and lepton sectors. We proceed by introducing the model, followed by an example of a flavor symmetry to showcase our results.

2. The model

The (stable) static spacetime background is:

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

where the extra coordinate y ranges between the two boundaries at $y = 0$ and $y = y_{\text{TeV}}$, and where $A(y)$ is the warp factor responsible for exponentially suppressing mass scales at different slices of the extra dimension. In the original RS scenario $A(y) = ky$, with k the curvature scale of the AdS_5 interval, while in general warped scenarios $A(y)$ is a more general (growing) function of y. The appeal of more complicated metrics lies on the possibility of having light KK resonances matter fields ($\sim$1 TeV), while keeping flavor and precision electroweak bounds at bay [8,11]. For simplicity we will take $A(y) = ky$, but with the assumption that the same general features arise in more complicated metric scenarios (this will be addressed in a longer companion paper). Assuming invariance under the usual SM gauge group, the 5D quark Lagrangian is

$$\mathcal{L}_q = \mathcal{L}_{kinetic} + M_{q_i}\bar{Q}_i Q_i + M_{u_i}\bar{U}_i U_i + M_{d_i}\bar{D}_i D_i + \left(Y^{5D}_{ij_u} H\bar{Q}_i U_j + \text{h.c.}\right) + \left(Y^{5D}_{ij_d} H\bar{Q}_i D_j + \text{h.c.}\right) \tag{2}$$

where Q_i, U_i and D_i are 5D quarks (doublets and singlets under $SU(2)$). In the lepton sector, we assume that Majorana mass terms are forbidden, and so the Lagrangian can be trivially obtained from the previous one by substituting Q_i by L_i, U_i by N_i and D_i by E_i, where L_i are lepton doublets, and N_i and E_i are neutrino and lepton singlets, respectively. The Higgs field H is a bulk scalar that can acquire a nontrivial vacuum expectation value (VEV) $v(y) = v_0 e^{aky}$, and is exponentially localized towards the TeV boundary, with delocalization controlled by the parameter a. Such nontrivial exponential VEVs appear naturally in warped backgrounds with simple scalar potentials and appropriate boundary conditions [8,12,13].

This extra dimensional scenario has two sources of flavor. One arises from the usual Yukawa couplings Y^u_{ij}, Y^d_{ij}, Y^e_{ij} and Y^ν_{ij} (dimensionless parameters defined in units of the curvature out the dimension-full 5D Yukawa couplings as $Y^{5D}_{ij} = \sqrt{k}Y_{ij}$). The other flavor parameters come from the fermion bulk mass terms, diagonal in flavor space, taken to be constant bulk terms written in units of the curvature k, i.e. $M_i = c_i k$ ($M_i = M_{q_i}, M_{u_i}, M_{d_i}, M_{L_i}, M_{\nu_i}, M_{e_i}$).

It has been observed before in [9] that, whenever the bulk Higgs localization parameter a is small enough in comparison with the c_i parameters (i.e., for the Higgs sufficiently delocalized from

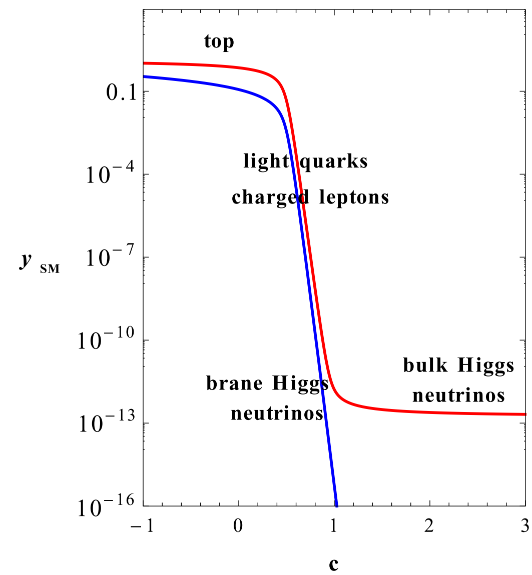

Fig. 1. Effective 4D Yukawa couplings for fermions as a function of the fermion bulk mass parameter c. For simplicity, we have taken the c-parameters for the doublet and the singlet to be equal.

the TeV brane), the 4D effective masses depend exponentially on a rather than on the c_i parameters. The effective 4D masses for all the SM fermions become

$$m_t \simeq v\tilde{Y}_{33} \quad c_{q_3}, c_{u_3} < 1/2 \tag{3}$$

$$(m_f)_{ij} \simeq v\epsilon^{(c_{L_i}-\frac{1}{2})}\epsilon^{(c_{R_j}-\frac{1}{2})}\tilde{Y}_{ij} \quad a > c_{L_i} + c_{R_j} \tag{4}$$

$$(m_\nu)_{ij} \simeq v\epsilon^{a-1}\tilde{Y}_{ij} \quad a < c_{l_i} + c_{\nu_j}, \tag{5}$$

where m_t is the top quark mass, $(m_f)_{ij}$ represents mass matrices for light quarks and charged leptons, and $(m_\nu)_{ij}$ is the Dirac neutrino mass matrix. Note that the couplings $\tilde{Y}_{ij}$ still retain a mild dependence on the c_i-parameters as follows $\tilde{Y}_{ij} = Y_{ij}\frac{\sqrt{2(a-1)(1-2c_{L_i})(1-2c_{R_j})}}{a-c_{L_i}-c_{R_j}}$, with Y_{ij} being the original 5D Yukawa couplings. The parameters $c_{L_i} \equiv c_{q_i}, c_{l_i}$ correspond to the $SU(2)$ doublets, and $c_{R_j} \equiv c_{u_j}, c_{d_j}, c_{e_j}, c_{\nu_j}$ are for the $SU(2)$ singlets. The warp factor ϵ defined by the background parameters as $\epsilon = e^{-ky_{\text{TeV}}} \sim 10^{-15}$ encapsulates the hierarchy between the UV (gravity) brane and the TeV (SM) brane. The c-parameters dependence on masses is shown in Fig. 1 for the case of a brane localized Higgs VEV ($a = 30$) and for a delocalized (bulk) Higgs VEV ($a = 1.9$), where the appearance of a plateau in the neutrino mass region reflects the insensitivity to the values of the c_i's in that limit.

Tension arises since, in order to generate viable neutrino masses from Eq. (5), one requires that $a \sim 1.80$–1.95. Values of $a < 2$ reintroduce some amount of tuning in the model and, for example, for $a = 1.95$, some ostensibly independent parameters of the 5D Higgs potential must be fixed to be equal to within about 1%. However this same tuning will also be responsible for generating a light enough Higgs mode compared to the KK scale [8]. In more general warped backgrounds this tension would easily disappear due to an enlarged parameter space, justifying the choice $a = 1.9$ throughout the rest of the paper.

We assume that all Yukawa matrices and fermion bulk masses from the 5D Lagrangian share the same symmetry structure, further broken by small terms i.e.

$$\mathbf{c}_f = \mathbf{c}^0_f + \delta\mathbf{c}_f \tag{6}$$

$$Y_Y = Y^0_Y + \delta Y_Y, \tag{7}$$

for all fermions of the model with $Y_Y = Y_u, Y_d, Y_e, Y_\nu$ and with $\mathbf{c}_f = \mathbf{c}_q, \mathbf{c}_u, \mathbf{c}_d, \mathbf{c}_l, \mathbf{c}_e, \mathbf{c}_\nu$. The matrices Y_Y^0 and $\mathbf{c}_f^0$ are flavor symmetric, while the small corrections $\delta \mathbf{c}_f$ and δY_Y quantifying the symmetry breaking do not have *a priori* any flavor structure.

From Eqs. (3)–(5), the fermion masses receive different corrections due to flavor violating terms:

$$m_t = m_t^0 + \delta m_t \quad c_{q_3}, c_{u_3} < 1/2 \tag{8}$$

$$(m_f)_{ij} = (m_f)_{ij}^0 \, \epsilon^{(\delta c_{L_i} + \delta c_{R_j})} \quad a > c_{L_i} + c_{R_j} \tag{9}$$

$$(m_\nu)_{ij} = (m_\nu)_{ij}^0 + \delta(m_\nu)_{ij} \quad a < c_{l_i} + c_{\nu_j}. \tag{10}$$

The exponential sensitivity of the fermion masses on the c-parameters is responsible for an exponential sensitivity of the symmetry breaking terms. Since $\epsilon \sim 10^{-15}$, the corrections to the mass matrices caused by these terms are of order $\sim 10^{-15(\delta c_i + \delta c_j)}$, which means that they could account for the observed hierarchies in the quark and charged lepton masses, as long as the symmetry breaking corrections δc_i remain between -0.1 and $+0.1$. The mixing angles are also exponentially sensitive to small symmetry breaking terms so that the mixing angles diagonalizing the mass matrices from the left will be $V_{ij} \sim \epsilon^{(\delta c_{L_i} - \delta c_{L_j})}$ for $(i < j)$ for quarks and charged leptons. Thus effects of the original symmetry are washed out in the quark and charged lepton sectors, while in the Dirac neutrino sector the sensitivity to the symmetry breaking is linear (i.e. small).

3. Flavor democracy

To qualify our assertions, we present an example of a symmetry with a democratic structure, broken by small perturbations. We impose a simple structure for all the flavor parameters of the model, namely one which remains invariant under family permutations [14]. This leads to a flavor structure where the 5D Yukawa couplings are invariant under $S_3 \times S_3$, while the 5D fermion bulk mass matrices are invariant under S_3 only. Under this symmetry, the 5D Yukawa couplings and to 5D fermion bulk mass matrices will be democratic, and will be parametrized as

$$Y_Y^D \propto \begin{pmatrix} 1 & 1 & 1 \\ 1 & 1 & 1 \\ 1 & 1 & 1 \end{pmatrix} \quad \text{and} \quad \mathbf{c}_f^D = \begin{pmatrix} A_f & B_f & B_f \\ B_f & A_f & B_f \\ B_f & B_f & A_f \end{pmatrix}. \tag{11}$$

Since the complete flavor structure is described by Eq. (11), we simultaneously diagonalize all matrices and obtain

$$Y_Y^0 = y_Y^0 \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix} \quad \text{and} \quad \mathbf{c}_f^0 = \begin{pmatrix} c_{f_1}^0 & 0 & 0 \\ 0 & c_{f_1}^0 & 0 \\ 0 & 0 & c_{f_3}^0 \end{pmatrix}, \tag{12}$$

where $y_Y^0 = y_u, y_d, y_e, y_\nu$ are complex Yukawa couplings in the up, down, charged lepton and neutrino Yukawa sectors. The matrix $\mathbf{c}_f^0$ is diagonal, with real entries $\mathbf{c}_{f_i}^0$. Democratic mass matrices produce two massless fermions and one massive one. The 0-th order CKM and PMNS matrices can be parametrized as

$$V_i^0 = \begin{pmatrix} \cos\theta_i^0 & \sin\theta_i^0 & 0 \\ -\sin\theta_i^0 & \cos\theta_i^0 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \tag{13}$$

where $i = \text{CKM, PMNS}$. Both matrices contain an angle θ_i^0 which is fixed not by the $S_3 \times S_3$ symmetry, but by the symmetry breaking terms, implemented in Eqs. (6) and (7). The small breaking of the symmetry is responsible for the (small) lift of the zero masses, yielding a viable neutrino spectrum, with one heavier and two lighter eigenstates with similar masses, and with a *normal hierarchy* ordering:

$$m_3 \sim \frac{m_\nu^0}{\sqrt{|\epsilon^{(2c_{l_3}-1)} - 1|}}, \tag{14}$$

$$m_1 \sim \delta Y_\nu m_\nu^0 \qquad m_2 \sim \delta Y_\nu m_\nu^0, \tag{15}$$

where $m_\nu^0 = v\epsilon^{a-1}$. We have kept the term in c_{l_3} since it can have an effect when $c_{l_3} \lesssim \frac{1}{2}$. The data constraining neutrino masses require $m_\nu^0 \sim (0.05–0.1)$ eV, which in turn fixes the size of the Higgs localization parameter a. The neutrino masses depend on δY_ν but not on δc_i's, which remain basically free (even the 0-th order $c_{\nu_i}^0$ are relatively free, as long as they satisfy $a < c_\nu + c_l$). For example, one finds that, in order to generate a viable neutrino mass hierarchy ratio $r = (|m_2|^2 - |m_1|^2)/(|m_3|^2 - |m_1|^2) \sim 0.03$, one needs $\delta Y_\nu \lesssim \sqrt{r} \sim 0.17$.

In the charged lepton and in the up-and down-quark sectors, the massless states are also lifted by the flavor symmetry breaking leaving a suppression proportional to δY. But here, in addition, the exponential dependence on the symmetry breaking parameters δc_{f_i} creates a hierarchy among all the masses. The third generation charged fermion masses are

$$m_t \sim m_t^0, \tag{16}$$

$$m_b \sim m_b^0 \epsilon^{\delta c_{e_3}}, \tag{17}$$

$$m_\tau \sim m_\tau^0 \epsilon^{(\delta c_{l_3} + \delta c_{e_3})}, \tag{18}$$

with the 0-th order masses[1] $m_t^0 = y_u^0 v$, $m_b^0 = y_d^0 v \epsilon^{c_{d_3}^0 - 1/2}$ and $m_\tau^0 = y_e^0 v \epsilon^{(c_{l_3}^0 + c_{e_3}^0 - 1)}$. The lighter fermion masses behave as

$$m_{f_2} \sim \delta Y_Y m_{f_1}^0 \epsilon^{(\delta c_{L_2} + \delta c_{R_2})}, \tag{19}$$

$$m_{f_1} \sim \delta Y_Y m_{f_1}^0 \epsilon^{(\delta c_{L_1} + \delta c_{R_1})}, \tag{20}$$

where $m_{f_2} \equiv m_c, m_s, m_\mu$, $m_{f_1} \equiv m_u, m_d, m_e$ and with $m_{f_1}^0 = v\epsilon^{(c_{L_1}^0 + c_{R_1}^0 - 1)}$. Note that in the flavor symmetric limit, the electron and muon, the down and strange quarks, and the up and charm quarks, are massless. The symmetry breaking produces nonzero masses proportional to the generic size of δY among these fermions,[2] with an added source of hierarchy due to the exponential dependence on the δc_i.

Thus, it is straightforward in this scenario to obtain phenomenologically viable quark and lepton masses by appropriately fixing the different δc_i within the constraint $|\delta c_i| \lesssim 0.1$. The hierarchies between fermion masses occur naturally and are under control since they depend exponentially on *small* numbers (hierarchical but not too hierarchical). Some masses and mixings still depend linearly on δY so that the typical size of these terms cannot be too small since, for instance, $\delta Y \lesssim m_t/m_c$ in the (extreme) limit where the charm quark c-parameters are top-like.

In this unified scenario, one can also generate the observed mixing angles in the CKM and PMNS matrices. The CKM entries become[3]

$$V_{us} \sim \epsilon^{(\delta c_{q_1} - \delta c_{q_2})}, \tag{21}$$

$$V_{cb} \sim \delta Y \epsilon^{(\delta c_{q_2} - \delta c_{q_3})} \tag{22}$$

[1] Assuming that $c_{q_3}^0, c_{u_3} < \frac{1}{2}$ and $c_{d_3}^0, c_{e_3}^0, c_{l_3}^0 > \frac{1}{2}$.

[2] The values of the light quark wavefunctions on the TeV brane are slightly higher than in usual RS scenarios to overcome the δY suppression. This induces stronger couplings with KK gluons and thus generically enhance potentially dangerous FCNC processes. These dangerous effects can be dealt with by invoking additional flavor constrains, such as enforcing exactly $c_{d_1} = c_{d_2}$ [15], or going to more general warped backgrounds where flavor bounds can be much milder [8].

[3] Strictly speaking, through out this paper, by V_{ij} we denote the absolute value of the respective matrix element.

Table 1
Zeroth order 5D fermion c-parameters. For simplicity, we also set all the 0-th order Yukawa coefficients to be universal $y_u^0 = y_d^0 = y_\nu^0 = y_e^0 = 4.4$ and the Higgs localization parameter to $a = 1.9$.

f	q	u	d	l	ν	e
$c^0_{f_1} (\equiv c^0_{f2})$	0.55	0.60	0.60	0.55	5.00	0.60
$c^0_{f_3}$	0.40	0.40	0.50	0.40	2.00	0.60

$$V_{ub} \sim \delta Y \epsilon^{(\delta c_{q_1} - \delta c_{q_3})}. \tag{23}$$

With respect to the 0-th order CKM matrix entries from Eq. (13), V_{us} receives a suppression exponentially sensitive to the difference between two small terms with respect to the CKM, which can easily reproduce the Cabibbo angle. The entries V_{cb} and V_{ub}, are lifted from the initial zero value, and acquire a double suppression, one parametric suppression, linear in the small Yukawa perturbations $\delta Y \sim 0.1$, caused by the broken S_3 symmetry, in addition to the exponential one. Assuming the usual ordering $\delta c_{q_1} > \delta c_{q2} > \delta c_{q_3}$, the ratio V_{cb}/V_{us} will then be of the correct order of magnitude for the typical size for δY and δc. The expected order of the ratio $V_{ub}/V_{cb} \sim V_{us}$ also remains realistic (up to order one factors not taken into account in the estimates). This last feature is generic in the usual RS scenarios.

The parametric dependence of the PMNS entries is different from that of the CKM matrix elements:

$$V_{e2} \sim \sin\theta^0_\nu \tag{24}$$

$$V_{e3} \sim \delta Y^\nu_{13} \sqrt{\left|\epsilon^{(2c_{l_3}-1)} - 1\right|} \tag{25}$$

$$V_{\mu 3} \sim \delta Y^\nu_{23} \sqrt{\left|\epsilon^{(2c_{l_3}-1)} - 1\right|}. \tag{26}$$

Contrary to the mixing angles in the quark sector, the value of V_{e2} is not suppressed and remains generically of $\mathcal{O}(1)$, fixed by the structure of the neutrino Yukawa flavor violating matrix δY^ν_{ij}. The entries V_{e3} and $V_{\mu 3}$ are lifted from zero, both depend on δY and, not only are they not further suppressed by exponential terms, but can actually be enhanced by exponential terms (as long as the approximation remains valid). In particular if $c_{l_3} \lesssim 1/2$, it is possible to lift the values of the mixing angles as shown in Eqs. (25) and (26). This feature is specific to the case $a < c_{l_3} + c_{\nu_j}$ and $c_{l_3} < 1/2$, and is not a generic feature in usual RS scenarios. More precise (and less compact) formulae will be presented elsewhere.

The observed mixing angles in the neutrino sector are most sensitive to the flavor structure of the neutrino Yukawa matrix δY^ν, but do not depend much on the charged lepton Yukawa matrix δY^l or on the δc_i. The bulk mass parameter of the third family lepton doublet should satisfy $c_{l_3} < 1/2$ in order to easily obtain larger mixing angles for small values of $\delta Y \sim 0.1$ (given that $V^{\exp}_{\mu 3} \sim 0.65$ and that $V^{\exp}_{e3} \sim 0.15$). This condition is very interesting as it is the same in the quark sector, where we require $c_{q_3} \lesssim 1/2$, to obtain a large top quark mass. This could be a hint of an additional family symmetry among the $SU(2)$ doublets of the third family.

Comparing the expressions for the V_{PMNS} mixing angles and the neutrino masses, the element δY_{23} must be larger that the rest of δY so that $\delta Y \sim \delta Y_{13} \sim \frac{\delta Y_{23}}{4}$.

In this scenario it is easy to find a set of 0-th order bulk parameters that reproduce the SM and that show the features described above. For example, a representative point for which the SM masses and mixing matrices are only a small perturbation away (of order 10% around an S_3 symmetric set of parameters), is shown in Table 1. Charged fermions results are not too sensitive to small deviations in the Yukawa couplings ($\lesssim$0.1), and once the δc_i's are fixed, the δY's can even be taken randomly as long as they remain at around 10%. One can then obtain generic charged fermion masses and mixings consistent with the SM, and any level of precision is possible by tuning these values.

In this particular setup, flavor violation bounds are not reduced, and the origin of CP phases is left unexplained since they appear from the flavor symmetry breaking terms. Nevertheless different flavor symmetries can alleviate the stringent bounds coming from FCNC's (specially in $K - \bar{K}$ mixing), if for example the CP violating phases in the down sector are suppressed. A thorough search of symmetries is beyond the scope of this paper, but looking for this type of effect is definitely a good guiding principle in the search of flavor symmetries in this context, as they would allow for lower KK masses. Also, as mentioned earlier, it is possible to work in scenarios where flavor violation and electroweak precision constraints are satisfied for much lower KK masses thanks to a modification of the background metric. Such a study in this same context of flavor symmetries is underway [16].

4. Conclusion

In this letter we have proposed a general framework in warped extra dimensions where the SM flavor structure is unified in all fermion sectors. In this scenario, all the fields are localized in the bulk and governed by the same flavor symmetry. Small breaking terms are introduced for the 5D bulk mass and Yukawa parameters. But while the quark and charged lepton sectors are dominated by the small flavor breaking in the bulk c-parameters, the (Dirac) neutrino sector is dominated by flavor symmetry breaking Yukawa couplings, and this result is robust. Neutrino flavor structure and charged lepton/quark flavor structure do not impact negatively upon each other. A permutation symmetry was studied to illustrate the idea, but other symmetries can also be invoked and explored within this framework. The main difference between the neutrino versus the charged lepton and quark sectors stems from allowing the Higgs field leak sufficiently out of the TeV brane so that the neutrino sector looses sensitivity on the 5D bulk mass parameters. This scenario can explain both masses of charged fermions and neutrinos, as well as reproduce the mixing matrices in both quark and lepton sectors. Within the flavor democracy example, we obtained simple expressions for the PMNS matrix elements, relating them in a simple way to the 5D parameters, Eqs. (24)–(26). We also obtained simple estimates of the neutrino masses as a function of these parameters, Eqs. (14)–(15), and found that they must obey a *normal* hierarchy. We found also that the localization of the third generation of lepton doublets $c_{l_3} > 1/2$ is quite constrained and finally the setup predicts that the ratio of the VPMNS mixing angles $V_{e3}/V_{\mu 3}$ is nontrivially given by the ratio of the flavor violating corrections to the Yukawa couplings $\delta Y_{13}/\delta Y_{23}$.

Acknowledgements

We thank NSERC for partial financial support under grant number SAP105354.

References

[1] L. Randall, R. Sundrum, Phys. Rev. Lett. 83 (1999) 3370;
L. Randall, R. Sundrum, Phys. Rev. Lett. 83 (1999) 4690.
[2] T. Gherghetta, A. Pomarol, Nucl. Phys. B 586 (2000) 141;
Y. Grossman, M. Neubert, Phys. Lett. B 474 (2000) 361.
[3] K. Agashe, A. Delgado, M.J. May, R. Sundrum, J. High Energy Phys. 0308 (2003) 050.
[4] G. Burdman, Phys. Rev. D 66 (2002) 076003;
G. Burdman, Phys. Lett. B 590 (2004) 86;
S.J. Huber, Nucl. Phys. B 666 (2003) 269;
M.S. Carena, A. Delgado, E. Ponton, T.M.P. Tait, C.E.M. Wagner, Phys. Rev. D 71 (2005) 015010;

M.S. Carena, A. Delgado, E. Ponton, T.M.P. Tait, C.E.M. Wagner, Phys. Rev. D 68 (2003) 035010.
[5] K. Agashe, G. Perez, A. Soni, Phys. Rev. D 71 (2005) 016002.
[6] C. Csaki, A. Falkowski, A. Weiler, J. High Energy Phys. 0809 (2008) 008;
K. Agashe, A. Azatov, L. Zhu, Phys. Rev. D 79 (2009) 056006.
[7] G. Cacciapaglia, C. Csaki, J. Galloway, G. Marandella, J. Terning, A. Weiler, J. High Energy Phys. 0804 (2008) 006;
A.L. Fitzpatrick, L. Randall, G. Perez, Phys. Rev. Lett. 100 (2008) 171604;
J. Santiago, J. High Energy Phys. 0812 (2008) 046;
M.C. Chen, K.T. Mahanthappa, F. Yu, Phys. Rev. D 81 (2010) 036004;
C. Csaki, A. Falkowski, A. Weiler, Phys. Rev. D 80 (2009) 016001.
[8] J.A. Cabrer, G. von Gersdorff, M. Quiros, New J. Phys. 12 (2010) 075012;
J.A. Cabrer, G. von Gersdorff, M. Quiros, Phys. Lett. B 697 (2011) 208;
J.A. Cabrer, G. von Gersdorff, M. Quiros, J. High Energy Phys. 1105 (2011) 083;
J.A. Cabrer, G. von Gersdorff, M. Quiros, J. High Energy Phys. 84 (2011) 035024;
J.A. Cabrer, G. von Gersdorff, M. Quiros, J. High Energy Phys. 1201 (2012) 033;
M. Quiros, arXiv:1311.2824 [hep-ph].
[9] K. Agashe, T. Okui, R. Sundrum, Phys. Rev. Lett. 102 (2009) 101801.
[10] B. Keren-Zur, P. Lodone, M. Nardecchia, D. Pappadopulo, R. Rattazzi, L. Vecchi, Nucl. Phys. B 867 (2013) 394, arXiv:1205.5803 [hep-ph].
[11] A. Carmona, E. Ponton, J. Santiago, J. High Energy Phys. 1110 (2011) 137;
S. Mert Aybat, J. Santiago, Phys. Rev. D 80 (2009) 035005;
A. Delgado, D. Diego, Phys. Rev. D 80 (2009) 024030.
[12] G. Cacciapaglia, C. Csaki, G. Marandella, J. Terning, J. High Energy Phys. 0702 (2007) 036, arXiv:hep-ph/0611358.
[13] P.R. Archer, J. High Energy Phys. 1209 (2012) 095, arXiv:1204.4730 [hep-ph].
[14] Among the earliest works on flavor permutation symmetries are:
H. Harari, H. Haut, J. Weyers, Phys. Lett. B 78 (1978) 459;
S. Pakvasa, H. Sugawara, Phys. Lett. B 73 (1978) 61;
S. Pakvasa, H. Sugawara, Phys. Lett. B 82 (1979) 105;
E. Derman, H.S. Tsao, Phys. Rev. D 20 (1979) 1207;
Y. Yamanaka, H. Sugawara, S. Pakvasa, Phys. Rev. D 25 (1982) 1895.
[15] J. Santiago, J. High Energy Phys. 0812 (2008) 046.
[16] Mariana Frank, Cherif Hamzaoui, Nima Pourtolami, Manuel Toharia, in progress.