arXiv:hep-ph/0010195v2 27 Oct 2000

Fermion Masses, Mixings and Proton Decay in a Randall-Sundrum Model

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Abstract

We consider a Randall-Sundrum model in which the Standard Model fermions and gauge bosons correspond to bulk fields. We show how the observed charged fermion masses and CKM mixings can be explained, without introducing hierarchical Yukawa couplings. We then study the impact on the mass scales associated with non-renormalizable operators responsible for proton decay, neutrino masses, and flavor changing neutral currents. Although mass scales as high as $10^{11} - 10^{12}$ GeV are in principle possible, dimensionless couplings of order 10^{-8} are still needed to adequately suppress proton decay. Large neutrino mixings seem to require new physics beyond the Standard Model.

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1 Introduction

Higher dimensional models of space-time with non-factorizable geometries have attracted much interest recently, especially because they may provide a solution to the gauge hierarchy problem. In the Randall-Sundrum approach [1] (see also [2]) the warp factor $\Omega = e^{-\pi kR}$ generates an exponential hierarchy between the effective fundamental mass scales on the two branes located at the orbifold fixed points in the extra dimension. If the brane separation is $kR \simeq 11$, the scale on one brane is of TeV-size, while the scale on the other brane is of order $M_{\rm Pl} \sim 10^{19}$ GeV. The AdS curvature k and the 5d Planck mass M_5 are both of order $M_{\rm Pl}$.

In the original proposal [1] the SM fields are all assumed to reside on the TeV-brane and only gravity propagates in the extra dimension. This setup is very economical in solving the hierarchy problem. However, one would naively expect that non-renormalizable operators in the 4d effective theory, now only TeV-scale suppressed, would lead to rapid proton decay, and unacceptably large neutrino masses and flavor changing neutral currents.

In the following we will explore to what extent this problem can be solved by moving the SM fermions away from the TeV-bane, without invoking ad hoc symmetries such as baryon and/or lepton number. The basic idea, already explored in ref. [3], exploits the fact that closer to the Planck-brane the effective fundamental scale is much larger than a TeV. However, the fermions cannot be moved arbitrarily far from the TeV-brane, in case their overlap with the Higgs field becomes too small to provide sufficiently large fermion masses. The Higgs field must reside near the TeV-brane if the model is expected to solve the hierarchy problem [4–6] (supersymmetry may help to avoid this constraint). The SM gauge bosons are assumed to be bulk fields to maintain gauge invariance. Taking into account constraints from the electroweak precision data [3–9], especially the W and Z boson mass ratio [5], the Kaluza Klein (KK) excitations of bulk gauge bosons and fermions are of order 10 TeV. Thus, some tuning of parameters is required to reproduce the measured W and Z boson masses.

Having the SM fermions in the bulk can also help explain the fermion mass hierarchy [3, 10, 11]. Heavy fermions are localized near the TeV-brane, where their overlap with Higgs fields is large, while the light fermions reside closer to the Planck-brane. We will demonstrate that this mechanism also generates acceptable quark mixings without invoking any flavor symmetry. The generation of fermion masses and quark mixings along these lines has

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already been explored in the case of flat extra dimensions [10, 11]. There, some additional physics has to be assumed in order to localize the fermions or provide an appropriate profile for the Higgs in the extra dimension. In the scenario considered here, these are automatically generated by the non-factorizable geometry.

In the following we present a set of "order one parameters", describing the localization of the bulk fermions, which successfully reproduces the measured fermion masses and quark mixings. We also study the impact on the mass scales associated with non-renormalizable operators responsible for proton decay, flavor changing neutral currents, and neutrino masses. Finally, we briefly discuss the issue of neutrino mixings.

2 5d fermions

We consider the non-factorizable metric [1]

$$ds^2 = e^{-2\sigma(y)}\eta_{\mu\nu}dx^\mu dx^\nu + dy^2, \tag{2.1}$$

where $\sigma(y) = k|y|$. The 4-dimensional metric is $\eta_{\mu\nu} = \text{diag}(-1, 1, 1, 1)$, k is the AdS curvature, and y denotes the fifth coordinate. The equation of motion for a fermion in curved space-time reads

$$E_a^M \gamma^a (\partial_M + \omega_M)\Psi + m_\Psi \Psi = 0, \tag{2.2}$$

where E_a^M is the fünfbein, γ^a are the Dirac matrices in flat space, and ω_M is the spin connection. The index M refers to objects in 5d curved space, the index a to those in tangent space. Under the Z_2 orbifold symmetry the fermions behave as $\Psi(-y)_\pm = \pm\gamma_5\Psi(y)_\pm$. Thus $\bar\Psi_\pm\Psi_\pm$ is odd and the Dirac mass term, which is also odd, can be parametrized as $m_\Psi = c\sigma'$. Here and in the following prime denotes differentiations with respect to the fifth coordinate. On the other hand, $\bar\Psi_\pm\Psi_\mp$ is even. Using the metric (2.1) one obtains for the left- and right-handed components of the Dirac spinor [3, 12]

$$[e^{2\sigma}\partial_\mu\partial^\mu + \partial_5^2 - \sigma'\partial_5 - M^2]e^{-2\sigma}\Psi_{L,R} = 0, \tag{2.3}$$

where $M^2 = c(c \pm 1)k^2 \mp c\sigma''$ and $\Psi_{L,R} = \pm\gamma_5\Psi_{L,R}$.

Decomposing the 5d fields as

$$\Psi(x^\mu, y) = \frac{1}{\sqrt{2\pi R}}\sum_{n=0}^\infty \Psi^{(n)}(x^\mu)f_n(y), \tag{2.4}$$

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one ends up with the zero mode wave function [3, 12]

$$f_0 = \frac{e^{(2-c)\sigma}}{N_0}, \qquad (2.5)$$

where

$$N_0^2 = \frac{e^{2\pi kR(1/2-c)} - 1}{2\pi kR(1/2 - c)}. \qquad (2.6)$$

Because of the orbifold symmetry, the zero mode of Ψ_+ (Ψ_-) is a left-handed (right-handed) Weyl spinor. For $c > 1/2$ ($c < 1/2$) the fermion is localized near the boundary at $y = 0$ ($y = \pi R$), i.e. at the Planck- (TeV-) brane.

3 Masses for bulk quarks and leptons

The zero modes of leptons and quarks acquire masses from their coupling to the Higgs field

$$\int d^4x \int dy \sqrt{-g} \lambda_{ij}^{(5)} H \bar{\Psi}_{i+} \Psi_{j-} \equiv \int d^4x \, m_{ij} \bar{\Psi}_{iR}^{(0)} \Psi_{jL}^{(0)} + \cdots, \qquad (3.7)$$

where $\lambda_{ij}^{(5)}$ are the 5d Yukawa couplings. The 4d Dirac masses are given by

$$m_{ij} = \int_{-\pi R}^{\pi R} \frac{dy}{2\pi R} \lambda_{ij}^{(5)} e^{-4\sigma} H(y) f_{0i+}(y) f_{0j-}(y). \qquad (3.8)$$

We assume that the Higgs profile has an exponential form which peaks at the TeV-brane

$$H(y) = H_0 e^{ak(|y|-\pi R)}. \qquad (3.9)$$

Using the known mass of the W-boson we can fix one parameter, which we take to be the amplitude H_0, in terms of the 5d weak gauge coupling $g^{(5)}$. The parameter a determines the width of the profile. If the profile satisfies the equations of motion we have $a = 4$ [13], but we will also consider other values.

Various constraints on the scenario with bulk gauge and fermion fields have been discussed in the literature [3–7, 9]. With bulk gauge fields for instance, the SM relationship between the gauge couplings and masses of the

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L	e	μ	τ
m [MeV]	0.5	106	1777
(A): c_L	0.681	0.591	0.537
(B): c_L	0.834	0.664	0.567
(A): M_4 [GeV]	4.02 10^9	1.98 10^7	1.21 10^6

Table 1: Lepton mass parameters in case of (A) left-right symmetry ($c_L = c_E$) and (B) delocalized right-handed fermions ($c_E = 1/2$).

Z and W bosons gets modified. The electroweak precision data then requires the lowest KK excitation of the gauge bosons to be heavier than about 10 TeV [5]. This bound is independent of the localization of the fermions in the extra dimension.

Furthermore, the KK excitations of the SM gauge bosons contribute to the electroweak precision observables. Their effect is relatively small if the SM fermions are localized towards the Planck-brane ($c > 1/2$). In this case KK masses around 1 TeV are sufficient for the corrections from the gauge boson excitations to be compatible with the experimental bounds [3, 6]. Therefore, the above constraint derived from the W and Z boson mass ratio is very important in this range of parameters. In the following we will assume that the mass of the first KK gauge boson is $m_1^{(G)} = 10$ TeV. The corresponding masses of the fermions are then in the range 10 to 13 TeV, for $0 < c < 1$.

Let us begin with the charged leptons. In the absence of Dirac masses for the neutrinos we can start with diagonal Yukawa couplings for the leptons. To avoid a new hierarchy in the 5d couplings, we assume $\lambda_{ii}^{(5)} = g^{(5)}$ (but will also discuss other possibilities below). We take $k = M_5 = \overline{M}_{\rm Pl}$, where $\overline{M}_{\rm Pl} = 2.44 \times 10^{18}$ GeV is the reduced Planck mass. From $m_1^{(G)} = 10$ TeV we determine the brane separation $kR = 10.83$ [8], and taking $a = 4$ we obtain $H_0 = 0.396 k/g^{(5)}$.

The lepton masses depend on the 5d mass parameters of the left- and right-handed fermions, c_L and c_E respectively, which enter (3.8). The experimentally known lepton masses do not fix these six c-parameters unambiguously. Hence we will concentrate on two special scenarios: (A) left-right symmetry, i.e. $c_L = c_E$; (B) delocalized right-handed fermions, i.e. $c_E = 1/2$. In table 1 we give the numerical values of the lepton mass parameters which reproduce the physical masses in both cases. The larger the 5d Dirac mass, i.e. c parameter of the fermion becomes, the greater is its localization at
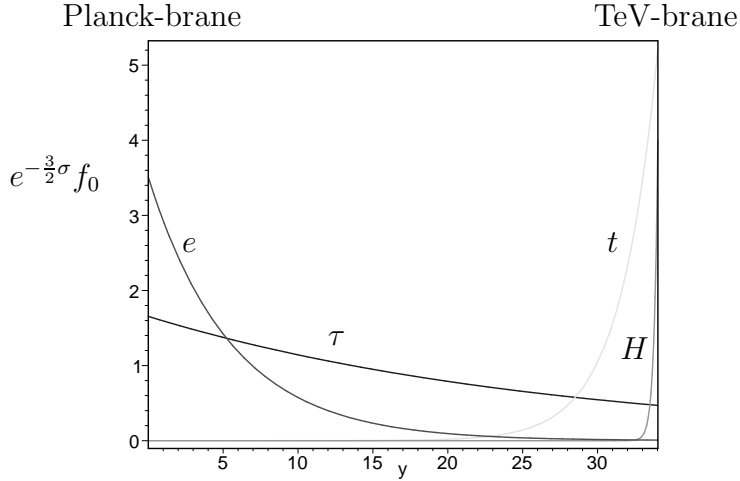
Figure 1: Localization of the electron ($c = 0.681$), tau ($c = 0.537$) and right-handed top quark ($c = 0.1$) zero modes in the fifth dimension. The Higgs profile H is given for $a = 4$ in units of $k^{3/2}$ and magnified by a factor of 10.

the Planck-brane. Its overlap with the Higgs profile at the TeV-brane is consequently less, which is reflected in a smaller 4d fermion mass after electroweak symmetry breaking. Our geometrical picture beautifully generates the charged lepton mass hierarchy by employing c-parameters of order unity. In figure 1 we present the wave functions of the electron and tau zero modes in the extra dimension, together with the Higgs profile. One observes that the electron, and to a lesser extent the tau, are localized near the Planck-brane. The factor of $e^{-\frac{3}{2}\sigma}$ compensates for the non-trivial measure induced by the AdS geometry.

The SM-like term (3.8) is not the sole contribution to the masses of the zero modes. The KK states of the fermions couple to the zero modes and induce corrections to (3.8) via tree-level and loop diagrams. We require these additional contributions to be small compared to (3.8). The tree-level contribution, arising due to a mixing of the zero modes and the excited fermion states, is of order [9]

$$\delta m_f^{(0)} \sim \lambda_{0i} \lambda_{ij} \lambda_{0j} \frac{v^3}{M_i M_j}, \tag{3.10}$$

where $M_{i,j}$ are the masses of the excited fermions, $v = 174$ GeV is the 4d Higgs vev and λ_{ij} are the Yukawa-type couplings between the fermion modes and the Higgs in the 4d effective theory. The corrections to fermion masses

from a loop involving a Higgs and KK fermion states has been estimated in [7],

$$\delta m_f^{(1)} \sim \frac{1}{16\pi^2} \lambda_{0i} \lambda_{ij} \lambda_{0j} v. \qquad (3.11)$$

Since the KK fermions are localized closer to the TeV-brane than the zero modes, their couplings to the Higgs boson is enhanced. Therefore, it is not clear that the fermion mass corrections involving these states are small, even though these states are rather heavy. Requiring that the corrections (3.10) and (3.11) are small compared to the fermion mass (3.7) provides additional constraints on the c-parameter of the fermions.

For the parameter sets (A) and (B) of table 1 we find that the additional contributions are indeed small, $\delta m_f^{(0)}/m_f \sim 0.03$ and $\delta m_f^{(1)}/m_f \sim 0.08$, irrespective of the lepton flavor. In this case the fermion masses can be reliably calculated from eq. (3.8).

If we increase the 5d Yukawa coupling, the fermions can be moved closer towards the Planck-brane. However, at the same time the fermion mass corrections (3.10) and (3.11) increase exponentially. We find that fine-tuning of the fermion mass contributions is only avoided for $\lambda^{(5)}/g^{(5)} \lesssim 5$. For the muon, for instance, this means that $c_L < 0.614$ ($c_L < 0.707$) for the left-right symmetric ($c_E = 1/2$) case. In this parameter range the constraints arising from the anomalous muon magnetic moment are satisfied as well [7].

In the case of quarks we have to take into account the additional complication of flavor mixing.[3] The mixing provides additional constraints on the 5d quark mass parameters. There are nine c-parameters in the quark sector, three each for the left-handed doublets c_Q, the right-handed u-quarks c_U, and the right-handed d-quarks c_D. The physical quantities we want to reproduce are the six quark masses and, in the absence of CP-violation, the three CKM mixing angles.

Rather then keeping the Yukawa couplings strictly equal to one, as in the case of leptons, we allow them to vary by a factor of two, i.e. $1/2 < |\lambda_{ij}^{(5)}|/g^{(5)} < 2$, which certainly introduces no new hierarchies. One parameter

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[3]For sake of simplicity we ignore CP-violation here. It is readily included by introducing complex Yukawa couplings for the quarks.

set which reproduces the physical quark masses and mixings [14] is given by

$$c_{Q1} = 0.72 \quad c_{D1} = 0.57 \quad c_{U1} = 0.63$$
$$c_{Q2} = 0.60 \quad c_{D2} = 0.57 \quad c_{U2} = 0.30$$
$$c_{Q3} = 0.35 \quad c_{D3} = 0.60 \quad c_{U3} = 0.10, \tag{3.12}$$

$$\frac{\lambda_D^{(5)}}{g^{(5)}} = \begin{pmatrix} 0.50 & -2.00 & -2.00 \\ 1.48 & -0.90 & 2.00 \\ 0.52 & -0.50 & 0.70 \end{pmatrix}, \quad \frac{\lambda_U^{(5)}}{g^{(5)}} = \begin{pmatrix} 0.80 & -1.90 & -2.00 \\ 1.23 & 1.20 & -1.04 \\ 1.85 & 1.66 & -0.80 \end{pmatrix}.$$

Using these numbers we obtain

$$m_u = 2.9 \text{ MeV}, \qquad m_c = 1.3 \text{ GeV},$$
$$m_d = 3.8 \text{ MeV}, \qquad m_b = 4.4 \text{ GeV},$$
$$m_s = 78 \text{ MeV}, \qquad m_t = 165 \text{ GeV},$$

$$|V_{\text{CKM}}| = \begin{pmatrix} 0.9744 & 0.2248 & 0.0045 \\ 0.2248 & 0.9736 & 0.0392 \\ 0.0045 & 0.0392 & 0.9992 \end{pmatrix}. \tag{3.13}$$

We have checked that for this parameter set the additional mass contributions from eqs. (3.10) and (3.11) are sub-leading. Like in the case of leptons an overall enhancement of the 5d Yukawa couplings by a factor of about five is tolerable. For larger Yukawa couplings, however, (3.10) and (3.11) become important.

Strictly speaking the quark masses given above are running masses at the cutoff scale of the effective 4d theory, which is in the TeV range. However, the effects of evolving the masses to the low energy regime could be absorbed in a redefinition of the 5d Yukawa mass matrices. The localization of the right-handed top-quark at the TeV-brane is shown in figure 1.

Note that the parameter set (3.12) is not uniquely determined by the experimental constraints. But it demonstrates that bulk fermions in the RS model can naturally explain with order one parameters not only the huge fermion mass hierarchy but also the quark mixings. Note that in the RS model a profile for fermions in the extra dimension is automatically induced by the non-factorizable geometry. As a result, this scenario is quite constrained, as will become clear in the discussion of the impact on non-renormalizable operators in the next section.

4 Dimension six operators and proton decay

We consider the following generic four-fermion operators which are relevant for proton decay and $K - \bar{K}$ mixing

$$\int d^4x \int dy \sqrt{-g} \frac{1}{M_5^3} \bar{\Psi}_i \Psi_j \bar{\Psi}_k \Psi_l \equiv \int d^4x \frac{1}{M_4^2} \bar{\Psi}_i^{(0)} \Psi_j^{(0)} \bar{\Psi}_k^{(0)} \Psi_l^{(0)} \qquad (4.14)$$

where the effective 4d mass scale is given by

$$\frac{1}{M_4^2} = \frac{1}{2\pi^2 k R^2 M_5^3} \frac{1}{N_0(c_i) N_0(c_j) N_0(c_k) N_0(c_l)} \frac{e^{(4-c_i-c_j-c_k-c_l)\pi k R} - 1}{4 - c_i - c_j - c_k - c_l}. \quad (4.15)$$

Let us first consider 4-fermion operators built from a single lepton flavor. For the left-right symmetric case and $M_5 = k$ the results for M_4 are given in the last line of table 1. One observes the rough relationship $M_4 \propto 1/m$. Furthermore we have $M_4 \propto \lambda^{(5)}/g^{(5)}$ and $M_4 \propto M_5^{3/2}$.

M_4 also depends on the width of the Higgs profile. The further the Higgs profile stretches out to the Planck brane (the smaller a gets), the closer the fermions can be moved to the Planck-brane, and the larger the suppression scale becomes. However, a has to be considerably smaller than four in order to change the given results. For instance, with $a = 2$, the effective M_4 is increased only by a factor of 2.5. Raising the suppression scale of a 4-fermion operator involving only electrons, for example, to 10^{16} GeV requires $c = 0.92$, which can be achieved only for $a \leq 1.4$. In this case the Higgs profile compensates the warp factor to a large extent. However, it is unclear if such a profile can be derived from a more fundamental scheme.

Let us now discuss some specific operators. Constraints from $K - \bar{K}$ mixing require that the dimension-six operator $(d\bar{s})^2/M_4^2$ is suppressed by $M_4 \gtrsim 10^6$ GeV. Using the c-parameters of (3.12) we obtain $M_4 = 5.5 \times 10^6$ GeV, in agreement with the constraint. This conclusion was also reached in ref. [3]. However, our estimate for the suppression scale is more realistic, since we have taken into account the quark mixings.

Concerning proton decay the most stringent constraints arise from the four-fermion operators $O_L = Q_1 Q_1 Q_2 L_3$ ($M_4 > 10^{15}$) GeV and $O_R = U_1^c U_2^c D_1^c E_3^c$ ($M_4 > 10^{12}$) GeV [15]. Using the c-parameters of (3.12) and case (B) of table 1 we obtain for these operators the effective suppression scales $M_4(O_L) = 7.7 \times 10^8$ GeV and $M_4(O_L) = 1.7 \times 10^6$ GeV, which are several orders of magnitude smaller than the experimental limits. If we enlarge

all the 5d Yukawa coupling by a factor of 5, the fermions can be moved further towards the Planck brane while still generating their desired masses. M_4 then increases by a factor of 5. Taking $M_5 = 10k$ instead of $M_5 = k$ would enlarge M_4 by another factor of about 30. If the two factors are combined, the effective suppression scales can be pushed to $M_4(O_L) = 1.2 \times 10^{11}$ GeV and $M_4(O_L) = 2.6 \times 10^8$ GeV, which is still four orders of magnitude below the constraint. Larger values of $\lambda^{(5)}$ and M_5 would further increase this result, but at the price of introducing new hierarchies. The same holds for a larger variation in the 5d Yukawa couplings. Moreover, for $\lambda^{(5)}/g^{(5)} \gtrsim 5$ the KK fermion contributions to the fermion masses (3.10) and (3.11) can no longer be ignored.

We conclude that by letting the SM fermions reside in the extra dimension considerable suppression of the non-renormalizable operators responsible for proton decay can be achieved. However, the effective suppression scale of these operators is still too small by at least four orders of magnitude. A small number of order 10^{-8} has to multiply the dimension-6 operators to satisfy the experimental constraints on the proton life time.

5   Dimension five neutrino masses

Majorana masses for left-handed neutrinos are generated by the dimension-five operator

$$\int d^4x \int dy \sqrt{-g} \frac{l_{ij}}{M_5^2} H^2 \Psi_{iL} C \Psi_{jL} \equiv \int d^4x \; m_{\nu ij} \Psi_{iL}^{(0)} C \Psi_{jL}^{(0)}, \qquad (5.16)$$

where l_{ij} are some dimensionless couplings constants, C is the charge conjugation operator and

$$m_{\nu ij} = \int_{-\pi R}^{\pi R} \frac{dy}{2\pi R} \frac{l_{ij}}{M_5^2} e^{-4\sigma(y)} H^2(y) f_{0i}(y) f_{0j}(y). \qquad (5.17)$$

Because of the SU(2) symmetry, the 5d Dirac mass parameters of the left-handed neutrinos and charged leptons are equal. No new parameters enter the game. For the left-right symmetric case (A) of table 1, taking $M_5 = k$ and $l_{ij} = \delta_{ij}$, i.e. ignoring neutrino mixing, we obtain

$$m_{\nu_e} = 39 \text{ keV}, \quad m_{\nu_\mu} = 8.8 \text{ MeV}, \quad m_{\nu_\tau} = 150 \text{ MeV}. \qquad (5.18)$$

These neutrino masses are well above the experimental limits.

The situation improves if the left-right symmetry is given up, and left-handed fields are moved closer to the Planck-brane. Let us consider case (B) where the right-handed fields are not localized at all ($c_E = 0.5$). Then the 5d mass parameters of the left-handed leptons read $c_e = 0.834$, $c_\mu = 0.664$ and $c_\tau = 0.567$. We then obtain the neutrino masses

$$m_{\nu_e} = 2.3 \text{ eV}, \quad m_{\nu_\mu} = 112 \text{ keV}, \quad m_{\nu_\tau} = 33 \text{ MeV}. \qquad (5.19)$$

While these neutrino masses may be compatible with collider experiments, they are problematic for neutrino oscillations and cosmology.

From eq. 5.17 one observes that $m_\nu \propto 1/M_5^2$. For $c_R = 0.5$ we also have $m_\nu \propto 1/(\lambda^{(5)})^2$. Assuming $M_5 = 10k$ (instead of $M_5 = k$) and $\lambda^{(5)}/g^{(5)} = 5$, we could reduce the neutrino masses by another factor of 2.5×10^3. However, the ν_τ mass is still in the keV region, and another factor of 10^5 is needed to bring it down to the range suggested by the atmospheric neutrino oscillation data. However, with the required large numbers, $M_5 = 100k$, $\lambda^{(5)}/g^{(5)} = 100$, new hierarchies arise in the model, and large corrections to fermions masses arise from eqs. (3.10) and (3.11). We conclude that some symmetry (e.g. lepton number) is needed to prevent large dimension five neutrino masses in the Randall-Sundrum model.

Finally, let us assume that with suitably large values of M_5 and $\lambda^{(5)}$ the neutrino masses of eq. 5.19 are indeed reduced by a factor of 10^8. Can the operator of eq. 5.16 explain the neutrino oscillation data? Including non-diagonal terms in (5.16) by setting all $l_{ij} = 1$ we obtain the following mass matrix in the $(\nu_e, \nu_\mu, \nu_\tau)$ basis

$$m_\nu = \begin{pmatrix} 2.28 & 505 & 8.66 \ 10^3 \\ 505 & 1.12 \ 10^5 & 1.92 \ 10^6 \\ 8.66 \ 10^3 & 1.92 \ 10^6 & 3.31 \ 10^7 \end{pmatrix} 10^{-8} \text{ eV}. \qquad (5.20)$$

As one might expect from the geometrical picture, the neutrino mass matrix is of nearest neighbor type, similar to those of the quarks. The corresponding mixings are in the few percent range, in conflict with the atmospheric neutrino data. If we relax the constraint on the couplings in (5.16) to $1/2 < |l_{ij}| < 2$ the situation improves to some extent. The ν_e–ν_μ (ν_μ–ν_τ) mixing angle can become as large as $\pi/5$ ($\pi/15$). But these values are still too low to explain the data. Moreover, we were not able to find a parameter set where both mixings are large at the same time, which seems favored by

the data. In order to explain the experimental results the field content of the model most likely has to be extended. The right-handed neutrino is an obvious choice (see, e.g. [12]), and we will explore this in a future publication.

Acknowledgement

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