Published by Institute of Physics Publishing for SISSA/ISAS
Received: *November 6, 2001*
Revised: *February 21, 2002*
Accepted: *March 5, 2002*

Signals from extra dimensions decoupled from the compactification scale

Francisco del Aguila and José Santiago

Departamento de Física Teórica y del Cosmos and
Centro Andaluz de Física de Partículas Elementales (CAFPE)
Universidad de Granada, E-18071 Granada, Spain
E-mail: faguila@ugr.es, jsantiag@ugr.es

Abstract: Multilocalization provides a simple way of decoupling the mass scale of new physics from the compactification scale of extra dimensions. It naturally appears, for example, when localization of fermion zero modes is used to explain the observed fermion spectrum, leaving low energy remnants of the geometrical origin of the fermion mass hierarchy. We study the phenomenology of the simplest five dimensional model with order one Yukawa couplings reproducing the standard fermion masses and mixing angles and with a light Kaluza-Klein quark $Q_{2/3}$ with observable new effects at large colliders.

Keywords: Field Theories in Higher Dimensions, Quark Masses and SM Parameters, Beyond Standard Model.

© SISSA/ISAS 2002

Contents

JHEP03(2002)010

1. Introduction

Theories with extra dimensions have received a great deal of attention during the last years [1, 2, 3]. They not only help to explain four dimensional puzzles but predict new physics at observable scales. Thus, the hierarchy of fermion masses and mixing angles can be related to the localization of the corresponding zero modes at different points in the extra dimensions [4]. These models reproduce the standard four dimensional fermion spectrum with order one Yukawa couplings, the small mass ratios and mixing angles resulting from the small overlapping of the corresponding wave functions in the extra dimensions [4]–[7].[1] However, the fermion splitting can induce large flavour changing neutral currents (FCNC) mediated by the Kaluza-Klein (KK) tower of gauge bosons [9], and then the experimental constraints on rare processes involving the first two families banish the masses of the first excited gauge bosons and the compactification scale M_c to high energy. A possibility of having remnants of the extra dimensions at low energy is to decouple the mass of the first KK modes from M_c. This can be done multilocalizing the zero mode, this is localizing it at two different points in the extra dimensions. Indeed, it has been proven in five dimensional models with localized gravity that the addition of mass terms modifies the localization properties of the fields, with the multilocalization of the first KK modes resulting in light four dimensional masses decoupled from the effective compactification scale [10]. This phenomenon, however, is general and independent of the gravitational background. As a matter of fact, it is a usual companion of nontrivial mass terms, and then can naturally appear in flavour models in extra dimensions.

[1] For another approach to the flavour problem in extra dimensions see [8].

In section 2 we study the simplest five dimensional model with multilocalized fermions. We assume a flat extra dimension compactified on S^1/Z_2 and step function masses, which can be thought as a limit of more realistic scalar backgrounds. In fact, the corresponding four dimensional effective lagrangian is the same as for more elaborate theories.[2] In an appendix we show that in models with split fermions compactified on S^1/Z_2 (like, for instance, the one recently proposed in [12]) there is also in general multilocalization, and then massless chiral fermions and light KK modes decoupled from M_c. We use these results in section 3 to construct a definite model with order one Yukawa couplings reproducing the standard quark masses and mixing angles and with an additional vector-like quark of charge 2/3, $Q_{2/3}$, near the electroweak scale and observable in forthcoming experiments, being directly produced and/or modifying the top couplings. Finally section 4 is devoted to phenomenological implications and conclusions.

2. Light Kaluza-Klein fermions in flat space

The appearance of new states parametrically lighter than the compactification scale is a consequence of multilocalization. It has been recently discussed in great detail in the case of warped compactifications for particles with spin smaller than two [10] (for the graviton it was first studied in [13] and further developed in [14].) It can be also present, however, in a flat background. We review in this section fermion multilocalization in the simplest possible context, a five dimensional model in flat space with the extra dimension compactified on the orbifold S^1/Z_2. This is a circle of radius R with the Z_2 identification $y \to -y$ or an interval $0 \leq y \leq \pi R$ with two boundaries, the orbifold fixed points.

The action of a spinor reads in this space (we use the "mostly minus" convention for the metric $(+,-,-,-,-)$ and $\gamma^4 = \mathrm{i}\gamma^5$)

$$S = \int \mathrm{d}^4x \int_0^{\pi R} \mathrm{d}y\ \bar{\Psi}[\mathrm{i}\gamma^N \partial_N - M(y)]\Psi\,, \qquad N = 0, \ldots, 4\,, \tag{2.1}$$

where the Dirac mass, which is odd under the action of Z_2, can be chosen to have a multi-kink structure in order to provide the desired multilocalization

$$M(y) = \begin{cases} M\,, & 0 \leq y \leq \pi a\,, \\ -M\,, & \pi a \leq y \leq \pi R\,, \end{cases} \tag{2.2}$$

with $0 \leq a \leq R$. Note that the mass M has to be real by hermiticity but there is no restriction on its sign. The kink-antikink shape for the mass term [15] is recovered in the limit $a \to R$. The KK reduction is performed in the usual way. We split the five dimensional vector-like fermion into its two chiralities $\Psi = \Psi_L + \Psi_R$ satisfying $\gamma^5 \Psi_{L,R} = \mp \Psi_{L,R}$, and expand in KK modes

$$\Psi_{L,R}(x,y) = \frac{1}{\sqrt{\pi R}} \sum_{n=0}^{\infty} f_n^{L,R}(y)\Psi_{L,R}^{(n)}(x)\,. \tag{2.3}$$

[2] This also applies to deconstructing models in four dimensions [11].

JHEP03(2002)010

Substituting them into the action, the decoupling of the quadratic terms follows from the relations

$$\int_0^{\pi R} \mathrm{d}y \, \frac{f_n^L f_m^L}{\pi R} = \int_0^{\pi R} \mathrm{d}y \, \frac{f_n^R f_m^R}{\pi R} = \delta_{nm} \,,$$
$$(\partial_y - M(y)) f_n^L = -m_n f_n^R \,, \qquad (-\partial_y - M(y)) f_n^R = -m_n f_n^L \,. \tag{2.4}$$

Then the corresponding four dimensional lagrangian describes a chiral zero mode plus a tower of vector-like fermions. The action of Z_2 on Ψ can be chosen to be $\Psi(-y) = \gamma^5 \Psi(y)$ or $-\gamma^5 \Psi(y)$. For the sake of concreteness we take the first choice, implying that the Right Handed (RH) component is even $\Psi_R(-y) = \Psi_R(y)$ and the Left Handed (LH) one odd $\Psi_L(-y) = -\Psi_L(y)$. The other assignment can be obtained from our results by just interchanging Ψ_L and Ψ_R and replacing M by $-M$. The coupled first order differential equations (2.4) imply the second order one

$$(-\partial_y^2 - M'(y) + M^2) f_n^R = m_n^2 f_n^R \,, \tag{2.5}$$

where the prime stands for ∂_y and

$$M'(y) = 2M[\delta(y) - \delta(y - \pi a) + \delta(y - \pi R)] \,. \tag{2.6}$$

f_n^L, which satisfies the same equation but with $M \to -M$, can be also obtained from f_n^R for $n \neq 0$ using the second equation in (2.4):

$$f_n^L = \frac{1}{m_n} (\partial_y + M(y)) f_n^R \,. \tag{2.7}$$

The solution of the Schrödinger equation (2.5) for f_n^R is obtained imposing the corresponding boundary conditions at $y = 0, \pi a$ and πR,

$$\begin{aligned} f_n^{R\prime}(0) &= -M f_n^R(0) \,, \\ f_n^R(\pi a - \epsilon) &= f_n^R(\pi a + \epsilon) \,, \\ f_n^{R\prime}(\pi a + \epsilon) - f_n^{R\prime}(\pi a - \epsilon) &= 2M f_n^R(\pi a) \,, \\ f_n^{R\prime}(\pi R) &= M f_n^R(\pi R) \,, \end{aligned} \tag{2.8}$$

where the limit $\epsilon \to 0$ is understood. The Z_2 projection leaves a chiral zero mode with even chirality, i.e. f_0^R. The odd boundary conditions, which imply the vanishing of the wave function at the orbifold fixed points, are not compatible with a massless mode. The zero mode wave function reads

$$f_0^R(y) = \begin{cases} A \exp[-M(y - \pi a)] \,, & 0 \leq y \leq \pi a \,, \\ A \exp[M(y - \pi a)] \,, & \pi a \leq y \leq \pi R \,, \end{cases} \tag{2.9}$$

where $A = [(2M\pi R)/(\exp[2M\pi a] + \exp[2M\pi(R - a)] - 2)]^{1/2}$ is a normalization constant. Thus it is exponentially localized at both orbifold fixed points for $M > 0$ and at $y = \pi a$ (intermediate brane) for $M < 0$. (The opposite would happen in the case of an even LH

JHEP03(2002)010

zero mode.) All the other KK fermions are vector-like, with the first massive mode $f_1^{L,R}$ having distinctive properties in the case of multilocalization. This occurs if the parameters in the potential, eqs. (2.5) and (2.6), satisfy

$$2M\pi a(R-a) > R\,. \tag{2.10}$$

Obviously, the case $a = R$ does not fulfil eq. (2.10) and there is no multilocalization for any value of the mass parameter. The phenomenology of this case [6, 7, 12, 16], which is also quite interesting, can be obtained from our results taking the continuous limit $a \to R$. In the case of multilocalization the first excited state is also exponentially localized

$$f_1^R(y) = \begin{cases} B\left(\mathrm{e}^{k_1 y} + \frac{k_1+M}{k_1-M}\mathrm{e}^{-k_1 y}\right), & 0 \leq y \leq \pi a\,, \\ C\left(\mathrm{e}^{-k_1(y-\pi R)} + \frac{k_1+M}{k_1-M}\mathrm{e}^{k_1(y-\pi R)}\right), & \pi a \leq y \leq \pi R\,, \end{cases} \tag{2.11}$$

and

$$f_1^L(y) = \begin{cases} D\left(\mathrm{e}^{k_1 y} - \mathrm{e}^{-k_1 y}\right), & 0 \leq y \leq \pi a\,, \\ E\left(\mathrm{e}^{-k_1(y-\pi R)} - \mathrm{e}^{k_1(y-\pi R)}\right), & \pi a \leq y \leq \pi R\,, \end{cases} \tag{2.12}$$

with B (D) and C (E) constants fixed by normalization and continuity conditions. Its mass $m_1^2 = M^2 - k_1^2$, with k_1 the solution of the eigenvalue equation

$$k - M - (k+M)\mathrm{e}^{-2k\pi R} + M\left[\mathrm{e}^{-2k\pi a} + \mathrm{e}^{-2k\pi(R-a)}\right] = 0, \tag{2.13}$$

can be $\ll M^2$. eq. (2.13) has always only one solution with $0 < k_1 < M$. The extreme case $m_1 \to 0$ $(m_1 \to M)$ corresponds to $2M\pi a(R-a) \gg R$ ($2M\pi a(R-a) \sim R$). The rest of the spectrum consists of oscillating states heavier than M and with a spacing of order $\sim 1/R$. In figure 1 we show the masses of the first KK fermions as function of the five dimensional mass M in units of $M_c = 1/R$ (we have fixed $a = R/2$ for illustration[3]). As can be observed, there is only one light KK vector-like fermion f_1 multilocalized for $MR > 2/\pi$, see eq. (2.10). In the absence of FCNC precision data typically require $M_c \gtrsim 4$ TeV, otherwise this limit can be as large as ~ 5000 TeV [9, 18]. This makes a priori the effects of the heavy states small compared to the contribution of the multilocalized one.

As has been recently discussed in [10] the appearance of multilocalization and the related light KK modes can be easily understood in terms of the shape of the potential in the corresponding Schrödinger equation. We review here the argument for completeness. It has been known for a long time that domain wall backgrounds create potential wells which can confine massless fermion zero modes in their world volume [19]. In the case that the potential has a double well shape with each well supporting a bound state (multidomain wall background), the resulting spectrum consists on two nearly degenerate modes with a mass splitting proportional to the quantum tunneling probability between wells. In the limit of large separation, the lightest state (which is always massless) corresponds to the even combination of the bound states for the separate wells, while the odd combination

[3] In this case the same model can be obtained with a kink-antikink mass compactifying on $S^1/(Z_2 \times Z_2')$ [17].

JHEP03(2002)010

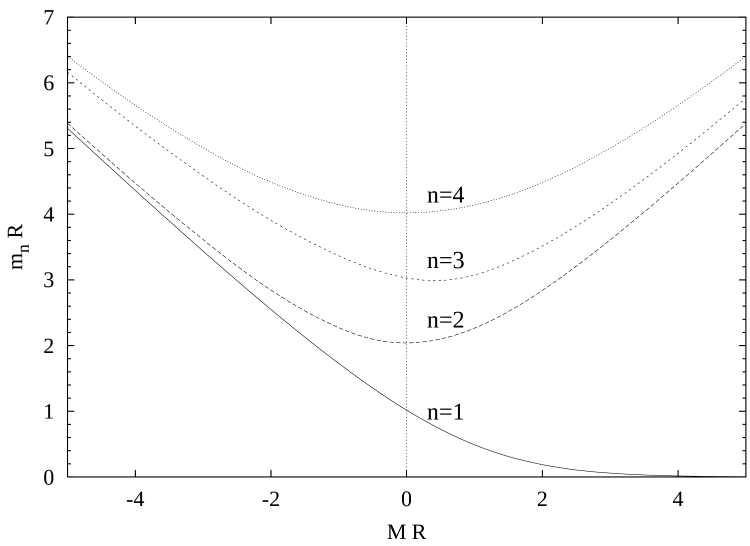

Figure 1: Values of the masses of the first KK modes as a function of the five dimensional mass in units of $1/R$. We have fixed $a = R/2$.

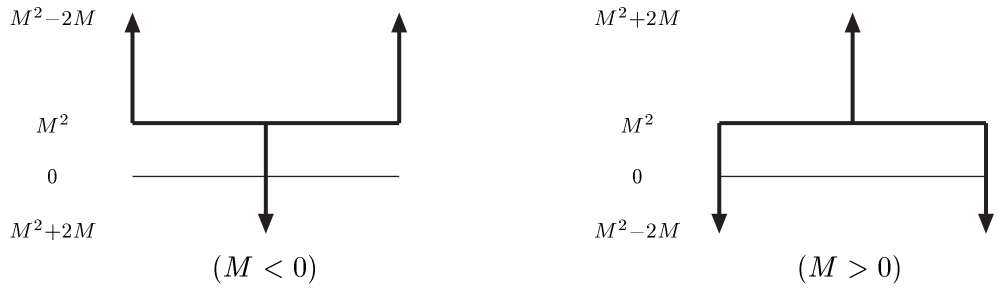

Figure 2: Potential $M^2 - M'(y)$ of the equivalent Schrödinger equation for a multikink mass term in arbitrary units (for $a = R/2$). On the left there is no multilocalization ($MR < 2/\pi$), in contrast with the potential on the right which does multilocalize ($MR > 2/\pi$).

has a mass proportional to the quantum tunneling probability. The absolute values of the corresponding wave functions only differ in the intermediate region where they are exponentially suppressed, what results in the exponentially small mass for the first excited mode. In figure 2 we draw the potential for the two signs of M in eq. (2.6): the case $M < 0$ (left) corresponds to the left region in figure 1 with no multilocalization, and the case $M > 0$ (right) to the right region in the same figure with the double well shape of the potential producing multilocalization for M large enough. The wave functions of the zero mode and the first KK excitation are shown in figure 3. Again the left (right) plot corresponds to the left (right) region in figure 1. It can be observed that in the latter case the difference between the absolute values of the wave functions are exponentially suppressed. Although we have discussed the particular case $a = R/2$, multilocalization will generically occur provided that the scalar background confines the fermion zero mode at separate points of the extra dimension. This happens for a wide range of the parameters,

JHEP03(2002)010

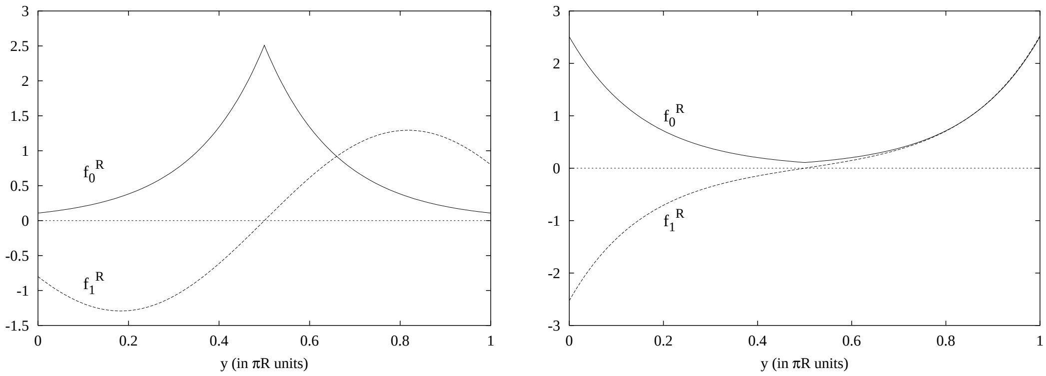

Figure 3: Profiles of the massless zero mode f_0^R and the first KK excitation f_1^R with no multilocalization, $MR = -2$ (left), and with multilocalization, $MR = 2$ (right).

location of the intermediate brane and strength of the potential, as shown in eq. (2.10).

In our case the potential consists of delta function wells and barriers because of the step function mass term we have considered. A possibly more realistic background would consist of hyperbolic-tangent shaped masses [15], with step functions being considered as a (thin brane) limit. From the discussion above it should be clear that in this more realistic case, provided that the scalar background has a three domain wall shape, multilocalization can be also present. As an example, we show in the Appendix that the alternative limit with the intermediate brane becoming fat and the boundary branes thin (which has been considered recently in [12], where the split fermion idea is realized on the orbifold S^1/Z_2) also presents multilocalization. Finally we would like to mention the importance that the orbifold has in our construction. As has been emphasized, multilocalization will generically occur provided that the scalar background giving mass to the five dimensional fermions has a multidomain wall structure. Orbifold models, which allow to obtain a chiral spectrum in five dimensions, naturally induce domain walls at both orbifold fixed points [15]. Thus, any scalar potential with an intermediate domain wall solution in the orbifold (what can be accomplished by the appropriate boundary sources [12]) automatically presents the multikink structure leading to multilocalization.

3. A model of flavour with light Kaluza-Klein quarks

In this section we construct a five dimensional model compactified on S^1/Z_2 with the standard fermion content, three $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ doublets q_i with LH components even and six singlets u_i, d_i with even RH components, which reduces below M_c to the Standard Model (SM) plus a light vector-like quark of charge 2/3 with sizeable mixing with the top quark. The mass terms, $M_i^{q,u,d}$, are order the compactification scale with step function shape and the Yukawa couplings, $\lambda_{ij}^{u,d(5)}$, order one. We shall make a detailed numerical discussion to show that there is no apparent fine tuning, and to emphasize its phenomenological relevance. Although the model is idealized, the four dimensional lagrangian is the same as for more realistic cases as already stated.

JHEP03(2002)010

The five dimensional action reads

$$S = \int \mathrm{d}^4x \int_0^{\pi R} \mathrm{d}y \left\{ \bar{q}_i \left[\mathrm{i}\gamma^N D_N - M_i^q(y) \right] q_i + (\mathrm{q} \to \mathrm{u}, \mathrm{d}) - \right.$$
$$\left. - \delta(y) \left[\lambda_{ij}^{u(5)} \bar{q}_i u_j \tilde{\phi} + \lambda_{ij}^{d(5)} \bar{q}_i d_j \phi + \mathrm{h.c.} \right] \right\}, \tag{3.1}$$

where D_N is the covariant derivative and to maintain the discussion simple we assume that the Higgs ϕ is at a fixed point, to be concrete at $y = 0$. We also assume that all the five dimensional masses M are generated from a unique scalar background and therefore that a is common to all of them. Performing the KK decomposition and integrating the extra dimension one obtains the four dimensional lagrangian

$$\mathcal{L} = \sum_{n=0}^{\infty} \left\{ \bar{q}_i^{(n)} [\mathrm{i}\gamma^\mu D_\mu - m_n^{q_i}] q_i^{(n)} + (\mathrm{q} \to \mathrm{u}, \mathrm{d}) \right\} -$$
$$- \sum_{n,m=0}^{\infty} \left\{ \lambda_{ij}^{u(nm)} \bar{q}_{Li}^{(n)} u_{Rj}^{(m)} \tilde{\phi} + \lambda_{ij}^{d(nm)} \bar{q}_{Li}^{(n)} d_{Rj}^{(m)} \phi + \mathrm{h.c.} \right\}, \tag{3.2}$$

where D_μ only includes the gauge boson zero modes and $m_0 = 0$. The effective four dimensional Yukawa couplings are

$$\lambda_{ij}^{u,d(nm)} = \frac{\lambda_{ij}^{u,d(5)}}{\pi R} f_{Ln}^{q_i} f_{Rm}^{u_j,d_j}, \tag{3.3}$$

where the wave functions f are evaluated at $y = 0$. The subscript L (R) for doublets (singlets) is skipped hereafter. For five dimensional masses $M \neq 0$ the exponential localization of the fermion zero modes can easily give the observed hierarchy of Yukawa couplings[4]

$$\lambda_{ij}^{u,d(00)} \sim \begin{pmatrix} \epsilon^4 & \epsilon^3 & \epsilon^2 \\ \epsilon^3 & \epsilon^2 & \epsilon \\ \epsilon^2 & \epsilon & 1 \end{pmatrix} \tag{3.4}$$

if $f_0^{q_i} \sim f_0^{u_i} \sim f_0^{d_i} \sim (\epsilon^2, \epsilon, 1)$. This matrix has two zero eigenvalues which become order ϵ^2 and ϵ^4 when the order one five dimensional Yukawa couplings are also included. To obtain such a hierarchy the third family must be strongly localized near the Higgs boundary and the first two families at $y = \pi a$, so that they are suppressed at the fixed point where the Higgs lives. As shown in figure 1 for M large enough (and $0 < a < R$) there is also multilocalization and a light KK vector-like fermion. In this scenario it seems natural to have t_R more strongly localized than t_L since the localization of the latter is the same as the one of its doublet counterpart b_L, which has a smaller Yukawa. This and the fact that with only one light vector-like quark we can account for large deviations of the SM top couplings make sufficient to have only t_R multilocalized.

[4] A similar texture for the Yukawa matrices is obtained in [20] (see also [21] for family symmetries giving the same mass matrices). The order of the different entries can be slightly changed varying the five dimensional Yukawa couplings.

JHEP03(2002)010

The effective lagrangian for the three light families is obtained integrating out the tower of KK modes in eq. (3.2). This has been done in detail in [16]. The corrections due to the KK fermions are proportional to the masses of the corresponding zero modes and to the inverse of the KK masses squared. Therefore only the top has large corrections and the main contribution is from the light multilocalized KK state. Indeed the largest top couplings to Z and $W^{\pm}$ in the mass eigenstate basis, $X_{tt}^{L,R}$[5] and W_{tb}^{L}, respectively, (W_{tb}^{R} has an extra suppression factor m_b/m_t) are to first order in $1/m_n^2$ [16]

$$X_{tt}^{L} = 1 - m_t^2 \sum_{k=1}^{3} |(U_R^u)_{kt}|^2 \sum_{n=1}^{\infty} \left(\frac{f_n^{u_k}}{m_n^{u_k} f_0^{u_k}} \right)^2 , \tag{3.5}$$

$$X_{tt}^{R} = m_t^2 \sum_{l,k,r=1}^{3} V_{tl} (U_L^q)^{\dagger}_{lk} \sum_{n=1}^{\infty} \left(\frac{f_n^{q_k}}{m_n^{q_k} f_0^{q_k}} \right)^2 (U_L^q)_{kr} V_{rt}^{\dagger} , \tag{3.6}$$

$$W_{tb}^{L} = V_{tb} - \frac{1}{2} m_t^2 \sum_{k=1}^{3} |(U_R^u)_{kt}|^2 \sum_{n=1}^{\infty} \left(\frac{f_n^{u_k}}{m_n^{u_k} f_0^{u_k}} \right)^2 . \tag{3.7}$$

We neglect corrections suppressed by $\frac{m_i V_{ib}}{m_t V_{tb}}$ with $i = u, c$, and the unitary matrices U diagonalize the zero mode mass submatrices

$$\begin{aligned} (U_L^q)^{\dagger}_{ik} \lambda_{kl}^{u(00)} \frac{v}{\sqrt{2}} (U_R^u)_{lj} &= V_{ij}^{\dagger} m_j^u , \qquad m_{1,2,3}^u = m_{u,c,t} , \\ (U_L^q)^{\dagger}_{ik} \lambda_{kl}^{d(00)} \frac{v}{\sqrt{2}} (U_R^u)_{lj} &= m_j^d , \qquad m_{1,2,3}^d = m_{d,s,b} , \end{aligned} \tag{3.8}$$

where v is the Higgs vev and, in the absence of KK corrections, $m_j^{u,d}$ are the mass eigenvalues and the unitary matrix V is the experimentally measured CKM matrix. On the other hand the corrected top mass reads

$$m_t^{\text{phys}} = m_t \left(1 - \frac{1}{2} m_t^2 \sum_{k=1}^{3} |(U_R^u)_{kt}|^2 \sum_{n=1}^{\infty} \left(\frac{f_n^{u_k}}{m_n^{u_k} f_0^{u_k}} \right)^2 \right) . \tag{3.9}$$

We are interested in large top mixing and a relatively light exotic quark consistent with present limits. Both requirements determine a and M_3^u once R is fixed, which is bounded to fulfil the constraints on FCNC due to the interchange of KK gauge bosons. Thus the model is essentially fixed if we further assume order one Yukawa couplings. For example, requiring $W_{tb}^L = 0.96$, $m_Q = 478\,\text{GeV}$ and $R = (85\,\text{TeV})^{-1}$, a fit to the data gives $a = 0.51R$,

$$\begin{aligned} M_i^q R &= (4.01, 2.26, -0.30) , \\ M_i^u R &= (-2.54, 0.75, 4.77) , \\ M_i^d R &= (-3.28, -2.88, -1.88) , \end{aligned} \tag{3.10}$$

and

$$\frac{\lambda_{ij}^{u(5)}}{\pi R} = \begin{pmatrix} 0.20 & -0.34 & 0.11 \\ -0.10 & 0.06 & -0.08 \\ 0.28 & 0.14 & 0.17 \end{pmatrix} , \qquad \frac{\lambda_{ij}^{d(5)}}{\pi R} = \begin{pmatrix} 0.35 & 0.10 & 0.46\,\mathrm{e}^{-2.77\,\mathrm{i}} \\ -0.17 & -0.65 & 0.10 \\ 0.25 & 0.13 & 0.25 \end{pmatrix} , \tag{3.11}$$

[5]They do not include the electromagnetic component, being equal to 1 and 0 in the SM, respectively.

JHEP03(2002)010

where the exponential dependence of the quark masses and mixing angles on the five dimensional masses make these very much constrained, while a wide range of variation is still allowed for the Yukawa couplings. With these values of the parameters the model reproduces the observed masses and mixing angles for the known quarks and predicts a vector-like quark of charge 2/3 with a mass in the current eigenstate basis $m_1^{u_3} = 462$ GeV. The other KK fermions have masses $m_n \gtrsim 50$ TeV. The observed CP violation is related to five dimensional Yukawa couplings which are in general complex, being however enough to take $\lambda_{13}^{d(5)}$ complex to obtain the measured CP violation. Indeed, integrating the KK modes [16], eqs. (3.5)–(3.9), we find (in GeV)

$$m_{u,c,t}^{\text{phys}} = (6 \times 10^{-3}, 2.6, 165)\,, \qquad m_{d,s,b}^{\text{phys}} = (11 \times 10^{-3}, 0.26, 6.6)\,, \tag{3.12}$$

where these masses are to be compared with the running $\overline{\text{MS}}$ masses evaluated at the scale m_t^{phys}. We have used the experimental values in [22], taking into account the running up to the scale m_t^{phys} [23],

$$\begin{aligned} m_u^{\text{phys}} &= 2.9 - 9.2 \text{ MeV}\,, & m_d^{\text{phys}} &= 5.5 - 16.5 \text{ MeV}\,, \\ m_c^{\text{phys}} &= 2.5 - 2.9 \text{ GeV}\,, & m_s^{\text{phys}} &= 140 - 310 \text{ MeV}\,, \\ m_t^{\text{phys}} &= 161 - 171 \text{ GeV}\,, & m_b^{\text{phys}} &= 6.2 - 6.8 \text{ GeV}\,. \end{aligned} \tag{3.13}$$

The corrected CKM matrix writes in the PDG phase convention [22]

$$W^L = \begin{pmatrix} 0.9750 & 0.222 & 0.003\ \mathrm{e}^{-1.11\,\mathrm{i}} \\ -0.222 - 0.00012\ \mathrm{e}^{1.11\,\mathrm{i}} & 0.9742 - 0.00004\ \mathrm{e}^{1.11\,\mathrm{i}} & 0.040 \\ 0.008 - 0.003\ \mathrm{e}^{1.11\,\mathrm{i}} & -0.038 - 0.0007\ \mathrm{e}^{1.11\,\mathrm{i}} & 0.96 \end{pmatrix}. \tag{3.14}$$

The amount of CP violation can be also given using the Jarlskog invariant [24]

$$|\text{Im}(W_{ub}^L W_{cs}^L W_{us}^{L*} W_{cb}^{L*})| = 2.4 \times 10^{-5}\,, \tag{3.15}$$

or the β angle [22]

$$\sin(2\beta) = \sin\left[2\text{arg}\left(-\frac{W_{cd}^L W_{cb}^{L*}}{W_{td}^L W_{tb}^{L*}}\right)\right] = 0.65\,, \tag{3.16}$$

which is in agreement with the most recent BaBar [25] and Belle [26] measurements

$$\sin(2\beta) = 0.59 \pm 0.14 \pm 0.05 \quad \text{and} \quad 0.99 \pm 0.14 \pm 0.06\,, \tag{3.17}$$

respectively. As can be observed, the masses and mixing angles in eqs. (3.12)–(3.14) give the correct experimental values, only differing appreciably from the minimal SM

$$W_{tb}^L = 0.96\,, \tag{3.18}$$

and then

$$X_{tt}^L = 0.93\,. \tag{3.19}$$

This reflects the lack of unitarity of the CKM matrix due to the top mixing with the lightest vector-like KK fermion, which is predicted to have an observable mass

$$m_Q = 478 \text{ GeV}\,. \tag{3.20}$$

JHEP03(2002)010

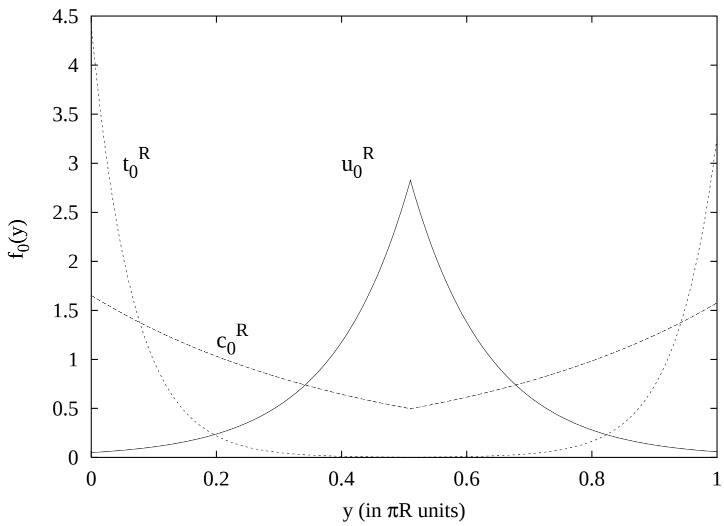

Figure 4: Profiles of the RH up quark zero modes for the five dimensional masses given in the text. Only t_R is multilocalized.

Several comments are in order. First, the required minimization to fit the experimental values fixing the masses and Yukawa couplings in eqs. (3.10) and (3.11) is essentially determined by the form of the four dimensional couplings in eq. (3.3) with their approximate geometrical form (3.4). Second, this structure is induced by the five dimensional masses through the (multi)localization of the zero modes, see for example figure 4. The quarks can be pushed to the four dimensional boundary taking the corresponding wave functions $f_{1,2,\dots} \to 0$ and $f_0 \to \sqrt{\pi R}$. Third, the fit has been done demanding $W^L_{tb} = 0.96$ and $m_Q = 478$ GeV. However, varying a and M^u_3 we can obtain any value of the top mixing and the exotic quark mass. In figure 5 we plot for arbitrary W^L_{tb} and m_Q the curves with a and M^u_3 fixed, respectively, our model depicted by a cross corresponding to $a = 0.51R$ and $M^u_3 R = 4.77$. Large departures from the SM in general imply large contributions to precision observables. In the absence of cancellations a large top mixing with a vector-like quark singlet is mainly constrained by $R_b^{\rm exp} = 0.21664 \pm 0.00068$ [22]. Using the results of ref. [27] we draw in figure 5 the 3 standard deviation excluded region for R_b (the SM prediction being at 1.4 standard deviations from the experimental value). In our model R_b is at 2.3 standard deviations, whereas the forward-backward asymmetry $A^b_{\rm FB}$ prediction is in slightly better agreement than the SM one and the new contributions to the oblique parameters S, T, U [28] remain within experimental errors [22].

4. Phenomenological implications and conclusions

We have discussed fermion multilocalization in a five dimensional model with the flat extra dimension compactified on S^1/Z_2 and step function masses. This phenomenon has been studied for particles up to spin 2 in warped backgrounds [10], but also happens in other backgrounds in the presence of domain wall masses. Multilocalization allows for a geometrical interpretation of the hierarchy of fermion masses and mixing angles and at

JHEP03(2002)010

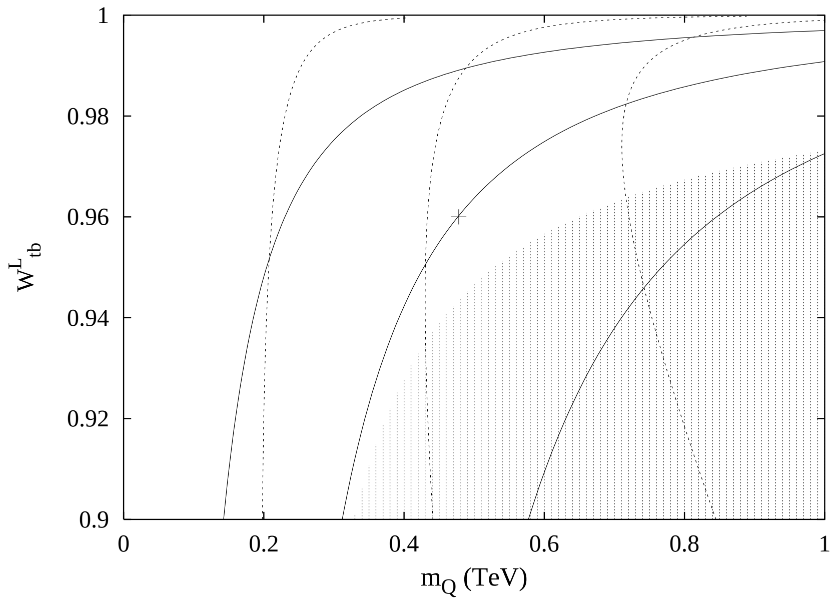

Figure 5: Top coupling W_{tb}^L and lightest vector-like quark mass m_Q as a function of the intermediate brane position a and the five dimensional mass M_3^u. Solid (dotted) lines stand for fixed $\frac{a}{R}$ $(M_3^u R)$ values, from left to right $0.53, 0.51, 0.49$ $(5.55, 4.85, 4.5)$. We take as in the text $R = (85 \quad \text{TeV})^{-1}$. The shadowed region corresponds to the 3 standard deviation exclusion region for R_b.

the same time for the decoupling of vector-like fermions from the compactification scale, making them observable.

In models with localized fermions at separate points in the extra dimensions there are a priori large FCNC, requiring compactification scales up to $M_c \gtrsim 5000$ TeV [9]. This bound is typically reduced for split fermion models with kink-antikink masses to $M_c \gtrsim 230$ TeV in the case of a boundary Higgs [12]. If these limits are not evaded by the specific model, and all KK excitations have masses order M_c or higher, no signal of the extra dimensions is expected at future colliders. However, this is not the case if there is multilocalization because the lightest KK modes decouple from M_c. The specific model we have worked out can accommodate a new vector-like quark of charge 2/3 with a mass $m_Q = 478$ GeV and a slightly lower compactification scale $M_c = 85$ TeV, evading all FCNC constraints.

Exotic quarks near the electroweak scale have been extensively studied in the past [29]. They are present in many grand unified models, like for instance E_6. Models in extra dimensions with the new vector-like quarks being the KK excitations of bulk fermions provide a natural realization of this possibility. In general the KK towers of vector-like fermions manifest at low energy modifying the fermion mixing [30], with the deviations from the SM predictions scaling with the masses of the SM quarks involved [16]. This which may be also indicated by experiment is naturally realized in these models, in contrast with former unified scenarios, what together with present experimental limits make the top, and then large colliders, the best place to look for these new effects. In the specific case we have considered Q can be directly observed, the reach of Tevatron being several hundreds GeV and of LHC few TeV [16, 31]. This vector-like quark also mixes with the

JHEP03(2002)010

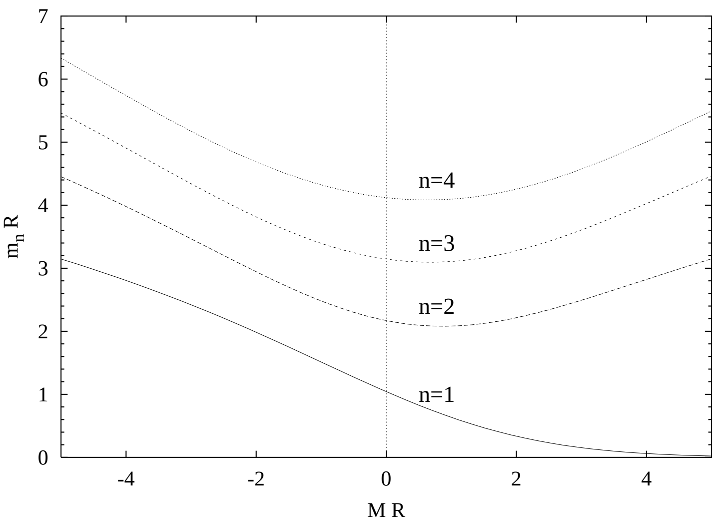

Figure 6: Values of the masses of the first KK modes for split fermions f_n as a function of the slope of the linear five dimensional mass.

top, modifying its couplings $W_{tb}^L = 0.96$ and $X_{tt}^L = 0.93$. Such a departure of the SM predictions, $W_{tb}^L \lesssim X_{tt}^L = 1$, cannot be established for W_{tb}^L at Tevatron Run II with an accumulated luminosity of 30 fb^{-1}, the expected accuracy being 7.6% [32] while it is at the edge for LHC with an expected precision of 5% [33]. Better prospects are for X_{tt}^L which will be measured with a precision of 2% at TESLA [34, 35]. In figure 5 we see that this model can reproduce any value of W_{tb}^L and m_Q varying the intermediate brane position and the five dimensional step function mass for a given compactification scale. However, banishing cancellations precision data restrict the top mixing and the exotic quark mass as shown in the figure.

As a final comment we would like to mention that we have concentrated on the quark sector, but multilocalization can be also used for leptons. A first attempt was made in [36] (see also [7]).

Acknowledgments

This work was supported in part by MCYT under contract FPA2000-1558, Junta de Andalucía group FQM 101 and the European Community's Human Potential Programme under contract HPRN-CT-2000-00149 Physics at Colliders. J.S. acknowledges MECD for finantial support under an FPU scholarship.

A. Split fermion multilocalization

In this appendix we show how multilocalization naturally appears in five dimensional models with split fermions and the flat extra dimension compactified on S^1/Z_2 [12]. Fermions are allowed to live in the bulk and have a linear odd mass term

$$M(y) = -\frac{M}{R}(y - \pi a)\,, \tag{A.1}$$

JHEP03(2002)010

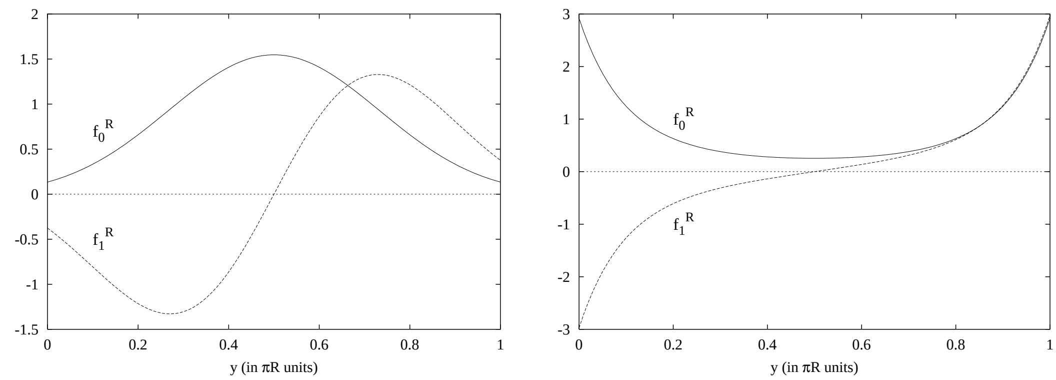

Figure 7: Profiles of the massless zero mode and the first KK excitation for the case of no multilocalization, $MR = -2$ (left), and multilocalization, $MR = 2$ (right).

with $0 \leq a \leq R$. A mass term of this kind can be obtained from a bulk scalar with appropriate source terms at the orbifold fixed points plus a step function mass term for $a \neq R/2$ [12]. For simplicity, we will consider the case $a = R/2$. As before M is constrained to be real by hermiticity but can have either sign. Let us consider, for definiteness, that the RH component of the fermion is even and the LH component odd. Again the opposite parity assignments can be obtained by just replacing RH by LH and M by $-M$. Equation (2.5) is still valid with the derivative of the mass term being now

$$M'(y) = -\frac{M}{R} + M\pi[\delta(y) + \delta(y - \pi R)]\,. \tag{A.2}$$

The corresponding Schrödinger equation in the interior of the interval is

$$\left[-\partial_y^2 + \frac{M}{R} + \frac{M^2}{R^2}(y - \pi R/2)^2\right] f_n^R = m_n^2 f_n^R\,. \tag{A.3}$$

In this case the potential is even under reflections about the middle point $y = \pi R/2$. Then we can solve in the region $\pi R/2 < y < \pi R$ and constrain the wave functions to be alternatively even and odd under this reflection. (Remember that we are solving for Ψ_R which is even under Z_2, and thus all wave functions are even at $y = \pi R$.) The boundary conditions read

$$f_{2n}^{R\,\prime}(\pi R/2) = f_{2n+1}^R(\pi R/2) = 0\,, \qquad f_n^{R\,\prime}(\pi R) = \frac{\pi M}{2} f_n^R(\pi R)\,, \qquad n = 0, 1, \ldots\,. \tag{A.4}$$

In order to solve the Schrödinger equation we first make the change of variables $t = \sqrt{\frac{M}{R}}(y - \pi R/2)$, which leads to the equation

$$[-\partial_t^2 + (1 - \lambda_n + t^2)] f_n^R = 0\,, \tag{A.5}$$

where we have defined the dimensionless variable $\lambda_n \equiv m_n^2 R/M$. We can now factorize the asymptotic form $f_n^R(y) = e^{-t^2/2} u_n(y)$, and make a further change of variables $z = t^2$. The resulting equation is a confluent hypergeometric equation

$$\left[z\partial_z^2 + \left(\frac{1}{2} - z\right)\partial_z - \frac{1}{4}(2 - \lambda_n)\right] u = 0\,, \tag{A.6}$$

JHEP03(2002)010

whose general solution can be written in terms of Kummer's functions $\mathcal{M}(a,b,z)$

$$u_n(z) = A\, \mathcal{M}\left(\frac{2-\lambda_n}{4}, \frac{1}{2}, z\right) + B\, z^{1/2}\, \mathcal{M}\left(1 - \frac{\lambda_n}{4}, \frac{3}{2}, z\right). \tag{A.7}$$

Inverting the changes of variables and applying the corresponding boundary conditions we find

$$f^R_{2n}(y) = A_{2n}\, e^{-\frac{M}{2R}(y-\pi R/2)^2}\, \mathcal{M}\left(\frac{2-\lambda_{2n}}{4}, \frac{1}{2}, \frac{M}{R}\left(y - \frac{\pi R}{2}\right)^2\right), \tag{A.8}$$

$$f^R_{2n+1}(y) = A_{2n+1}\left(y + \frac{\pi R}{2}\right) e^{-\frac{M}{2R}\left(y-\frac{\pi R}{2}\right)^2}\, \mathcal{M}\left(1 - \frac{\lambda_{2n+1}}{4}, \frac{3}{2}, \frac{M}{R}\left(y - \frac{\pi R}{2}\right)^2\right), \tag{A.9}$$

where the masses $m_n^2 = \lambda_n M/R$ are obtained solving the eigenvalue equations for λ_n:

$$\left(1 - \frac{\lambda_{2n}}{2}\right) \mathcal{M}\left(\frac{3}{2} - \frac{\lambda_{2n}}{4}, \frac{3}{2}, \frac{\pi^2 MR}{4}\right) - \mathcal{M}\left(\frac{1}{2} - \frac{\lambda_{2n}}{4}, \frac{1}{2}, \frac{\pi^2 MR}{4}\right) = 0\,, \tag{A.10}$$

for the even modes, and

$$\begin{aligned} &\left(1 - \frac{\pi^2 MR}{2}\right) \mathcal{M}\left(1 - \frac{\lambda_{2n+1}}{4}, \frac{3}{2}, \frac{\pi^2 MR}{4}\right) + \\ &+\frac{\pi^2 MR}{3}\left(1 - \frac{\lambda_{2n+1}}{4}\right) \mathcal{M}\left(2 - \frac{\lambda_{2n+1}}{4}, \frac{5}{2}, \frac{\pi^2 MR}{4}\right) = 0\,, \end{aligned} \tag{A.11}$$

for the odd ones.

To gain some intuition we start considering the massless zero mode whose wave function can be written in terms of elementary functions. Using the property $\mathcal{M}(a,a,z) = e^z$ we find that $\lambda_0 = 0$ is a solution of the eigenvalue equation for the even modes. The corresponding zero mode has a gaussian profile

$$f^R_0(y) = A_0\, e^{\frac{M}{2R}(y-\pi R/2)^2}\,, \tag{A.12}$$

being localized at $y = \pi R/2$ for $M < 0$ and at both fixed points with exponential suppression at the midpoint if $M > 0$. In the latter case the zero mode is multilocalized and we expect that the first KK mode be anomalously light. In figure 6 we plot the masses of the first few KK modes as a function of the slope of the five dimensional mass. The phenomenon of multilocalization and the corresponding light KK excitation is apparent for the appropriate M values. We also show in figure 7 the wave function profiles for the zero and first modes for the two M signs: $MR = -2$ with no multilocalization on the left; and $MR = 2$ with multilocalization on the right. There is a close resemblance of this spectrum with the one found in section 2 for the case of a multikink mass term. This could have been anticipated from the shape of the potential with also two attractive delta functions at the orbifold fixed points.

As a final remark we would like to comment that in the flavour models with split fermions proposed in the literature the appropriate signs for the mass terms have been chosen (positive for even LH fields and negative for even RH ones) so as to have gaussian localization around one single point and not multilocalization. However, we find no reason for a generic choice of sign excluding the possibility of multilocalization.

JHEP03(2002)010

References

[1] I. Antoniadis, *A possible new dimension at a few TeV*, *Phys. Lett.* **B 246** (1990) 377;
I. Antoniadis and K. Benakli, *Limits on extra dimensions in orbifold compactifications of superstrings*, *Phys. Lett.* **B 326** (1994) 69 [hep-th/9310151];
I. Antoniadis, K. Benakli and M. Quirós, *Production of Kaluza-Klein states at future colliders*, *Phys. Lett.* **B 331** (1994) 313 [hep-ph/9403290];
I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G.R. Dvali, *New dimensions at a millimeter to a fermi and superstrings at a TeV*, *Phys. Lett.* **B 436** (1998) 257 [hep-ph/9804398];
K.R. Dienes, E. Dudas and T. Gherghetta, *Extra spacetime dimensions and unification*, *Phys. Lett.* **B 436** (1998) 55 [hep-ph/9803466];
K.R. Dienes, E. Dudas and T. Gherghetta, *Grand unification at intermediate mass scales through extra dimensions*, *Nucl. Phys.* **B 537** (1999) 47 [hep-ph/9806292].

[2] N. Arkani-Hamed, S. Dimopoulos and G.R. Dvali, *The hierarchy problem and new dimensions at a millimeter*, *Phys. Lett.* **B 429** (1998) 263 [hep-ph/9803315]; *Phenomenology, astrophysics and cosmology of theories with sub-millimeter dimensions and TeV scale quantum gravity*, *Phys. Rev.* **D 59** (1999) 086004 [hep-ph/9807344].

[3] L. Randall and R. Sundrum, *A large mass hierarchy from a small extra dimension*, *Phys. Rev. Lett.* **83** (1999) 3370 [hep-ph/9905221]; *An alternative to compactification*, *Phys. Rev. Lett.* **83** (1999) 4690 [hep-th/9906064].

[4] N. Arkani-Hamed and M. Schmaltz, *Hierarchies without symmetries from extra dimensions*, *Phys. Rev.* **D 61** (2000) 033005 [hep-ph/9903417].

[5] E.A. Mirabelli and M. Schmaltz, *Yukawa hierarchies from split fermions in extra dimensions*, *Phys. Rev.* **D 61** (2000) 113011 [hep-ph/9912265];
G.R. Dvali and M.A. Shifman, *Families as neighbors in extra dimension*, *Phys. Lett.* **B 475** (2000) 295 [hep-ph/0001072];
G.C. Branco, A. de Gouvea and M.N. Rebelo, *Split fermions in extra dimensions and CP-violation*, *Phys. Lett.* **B 506** (2001) 115 [hep-ph/0012289];
G. Barenboim, G.C. Branco, A. de Gouvea and M.N. Rebelo, *Neutrino masses and lepton flavor violation in thick brane scenarios*, *Phys. Rev.* **D 64** (2001) 073005 [hep-ph/0104312].

[6] Y. Grossman and M. Neubert, *Neutrino masses and mixings in non-factorizable geometry*, *Phys. Lett.* **B 474** (2000) 361 [hep-ph/9912408];
T. Gherghetta and A. Pomarol, *Bulk fields and supersymmetry in a slice of AdS*, *Nucl. Phys.* **B 586** (2000) 141 [hep-ph/0003129];
S.J. Huber and Q. Shafi, *Fermion masses, mixings and proton decay in a Randall-Sundrum model*, *Phys. Lett.* **B 498** (2001) 256 [hep-ph/0010195].

[7] S.J. Huber and Q. Shafi, *Neutrino oscillations and rare processes in models with a small extra dimension*, *Phys. Lett.* **B 512** (2001) 365 [hep-ph/0104293].

[8] L. Hall, J. March-Russell, T. Okui and D.R. Smith, *Towards a theory of flavor from orbifold guts*, hep-ph/0108161.

[9] C.D. Carone, *Electroweak constraints on extended models with extra dimensions*, *Phys. Rev.* **D 61** (2000) 015008 [hep-ph/9907362];
A. Delgado, A. Pomarol and M. Quirós, *Electroweak and flavor physics in extensions of the standard model with large extra dimensions*, *J. High Energy Phys.* **01** (2000) 030 [hep-ph/9911252].

JHEP03(2002)010

[10] I.I. Kogan, S. Mouslopoulos, A. Papazoglou and G.G. Ross, *Multi-localization in multi-brane worlds*, *Nucl. Phys.* **B 615** (2001) 191 [hep-ph/0107307].

[11] N. Arkani-Hamed, A.G. Cohen and H. Georgi, *(De)constructing dimensions*, *Phys. Rev. Lett.* **86** (2001) 4757 [hep-th/0104005];
C.T. Hill, S. Pokorski and J. Wang, *Gauge invariant effective lagrangian for Kaluza-Klein modes*, *Phys. Rev.* **D 64** (2001) 105005 [hep-th/0104035].

[12] D.E. Kaplan and T.M.P. Tait, *New tools for fermion masses from extra dimensions*, *J. High Energy Phys.* **11** (2001) 051 [hep-ph/0110126].

[13] I.I. Kogan, S. Mouslopoulos, A. Papazoglou, G.G. Ross and J. Santiago, *A three three-brane universe: new phenomenology for the new millennium?*, *Nucl. Phys.* **B 584** (2000) 313 [hep-ph/9912552].

[14] S. Mouslopoulos and A. Papazoglou, *'$+-+$' brane model phenomenology*, *J. High Energy Phys.* **11** (2000) 018 [hep-ph/0003207];
I.I. Kogan, S. Mouslopoulos and A. Papazoglou, *A new bigravity model with exclusively positive branes*, *Phys. Lett.* **B 501** (2001) 140 [hep-th/0011141];
I.I. Kogan, S. Mouslopoulos, A. Papazoglou and G.G. Ross, *Multigravity in six dimensions: generating bounces with flat positive tension branes*, *Phys. Rev.* **D 64** (2001) 124014 [hep-th/0107086];
I.I. Kogan and G.G. Ross, *Brane universe and multigravity: modification of gravity at large and small distances*, *Phys. Lett.* **B 485** (2000) 255 [hep-th/0003074];
I.I. Kogan, S. Mouslopoulos, A. Papazoglou and G.G. Ross, *Multi-brane worlds and modification of gravity at large scales*, *Nucl. Phys.* **B 595** (2001) 225 [hep-th/0006030].

[15] H. Georgi, A.K. Grant and G. Hailu, *Chiral fermions, orbifolds, scalars and fat branes*, *Phys. Rev.* **D 63** (2001) 064027 [hep-ph/0007350].

[16] F. del Aguila and J. Santiago, *Universality limits on bulk fermions*, *Phys. Lett.* **B 493** (2000) 175 [hep-ph/0008143].

[17] R. Barbieri, L.J. Hall and Y. Nomura, *A constrained standard model from a compact extra dimension*, *Phys. Rev.* **D 63** (2001) 105007 [hep-ph/0011311].

[18] T.G. Rizzo and J.D. Wells, *Electroweak precision measurements and collider probes of the standard model with large extra dimensions*, *Phys. Rev.* **D 61** (2000) 016007 [hep-ph/9906234].

[19] R. Jackiw and C. Rebbi, *Solitons with fermion number* 1/2, *Phys. Rev.* **D 13** (1976) 3398;
V.A. Rubakov and M.E. Shaposhnikov, *Do we live inside a domain wall?*, *Phys. Lett.* **B 125** (1983) 136.

[20] R.G. Roberts, A. Romanino, G.G. Ross and L. Velasco-Sevilla, *Precision test of a fermion mass texture*, *Nucl. Phys.* **B 615** (2001) 358 [hep-ph/0104088].

[21] S.F. King and G.G. Ross, *Fermion masses and mixing angles from* SU(3) *family symmetry*, *Phys. Lett.* **B 520** (2001) 243 [hep-ph/0108112].

[22] Particle Data Group collaboration, D.E. Groom et al., *Review of particle physics*, *Eur. Phys. J.* **C 15** (2000) 1.

[23] V.D. Barger, M.S. Berger, T. Han and M. Zralek, *Test of the Dimopoulos-Hall-Raby ansatz for fermion mass matrices*, *Phys. Rev. Lett.* **68** (1992) 3394 [hep-ph/9203220];
G.W. Anderson, S. Raby, S. Dimopoulos and L.J. Hall, *Precise predictions for m_t, V_{cb} and* $\tan\beta$, *Phys. Rev.* **D 47** (1993) 3702 [hep-ph/9209250].

JHEP03(2002)010

[24] C. Jarlskog, *Commutator of the quark mass matrices in the standard electroweak model and a measure of maximal CP-violation*, *Phys. Rev. Lett.* **55** (1985) 1039; *A basis independent formulation of the connection between quark mass matrices, CP-violation and experiment*, *Z. Physik* **C 29** (1985) 491.

[25] BABAR collaboration, B. Aubert et al., *Observation of CP-violation in the B^0 meson system*, *Phys. Rev. Lett.* **87** (2001) 091801 [hep-ex/0107013].

[26] BELLE collaboration, K. Abe et al., *Observation of large CP-violation in the neutral B meson system*, *Phys. Rev. Lett.* **87** (2001) 091802 [hep-ex/0107061].

[27] P. Bamert, C.P. Burgess, J.M. Cline, D. London and E. Nardi, *R_b and new physics: a comprehensive analysis*, *Phys. Rev.* **D 54** (1996) 4275 [hep-ph/9602438].

[28] L. Lavoura and J.P. Silva, *The oblique corrections from vector-like singlet and doublet quarks*, *Phys. Rev.* **D 47** (1993) 2046.

[29] F. del Aguila and M.J. Bowick, *The possibility of new fermions with $\Delta I = 0$ mass*, *Nucl. Phys.* **B 224** (1983) 107;
P. Fishbane, S. Meshkov and P. Ramond, *Standard model constraints on fermions*, *Phys. Lett.* **B 134** (1984) 81;
F. del Aguila and J. Cortés, *A new model of weak CP-violation*, *Phys. Lett.* **B 156** (1985) 243;
G.C. Branco and L. Lavoura, *On the addition of vector like quarks to the standard model*, *Nucl. Phys.* **B 278** (1986) 738;
P. Langacker and D. London, *Mixing between ordinary and exotic fermions*, *Phys. Rev.* **D 38** (1988) 886;
R. Barbieri and L.J. Hall, *An exotic top quark*, *Nucl. Phys.* **B 319** (1989) 1;
J.L. Hewett and T.G. Rizzo, *Low-energy phenomenology of superstring inspired E_6 models*, *Phys. Rept.* **183** (1989) 193;
Y. Nir and D.J. Silverman, *Z mediated flavor changing neutral currents and their implications for CP asymmetries in B^0 decays*, *Phys. Rev.* **D 42** (1990) 1477;
E. Nardi, E. Roulet and D. Tommasini, *A simultaneous analysis of Z' and new fermions effects. Global constraints in E_6 and* SO(10) *models*, *Phys. Rev.* **D 46** (1992) 3040;
V.D. Barger, M.S. Berger and R.J.N. Phillips, *Quark singlets: implications and constraints*, *Phys. Rev.* **D 52** (1995) 1663 [hep-ph/9503204];
P.H. Frampton, P.Q. Hung and M. Sher, *Quarks and leptons beyond the third generation*, *Phys. Rept.* **330** (2000) 263 [hep-ph/9903387];
K. Higuchi and K. Yamamoto, *Quark mixings and flavor changing interactions with singlet quarks*, *Phys. Rev.* **D 62** (2000) 073005 [hep-ph/0004065];
Y. Liao and X. Li, *Non-unitarity of CKM matrix from the vector singlet quark mixing and neutron electric dipole moment*, *Phys. Lett.* **B 503** (2001) 301 [hep-ph/0005063].

[30] F. del Aguila, M. Pérez-Victoria and J. Santiago, *Effective description of quark mixing*, *Phys. Lett.* **B 492** (2000) 98 [hep-ph/0007160]; *Observable contributions of new exotic quarks to quark mixing*, *J. High Energy Phys.* **09** (2000) 011 [hep-ph/0007316].

[31] F. del Aguila, L. Ametller, G.L. Kane and J. Vidal, *Vector like fermion and standard Higgs production at hadron colliders*, *Nucl. Phys.* **B 334** (1990) 1;
T.G. Rizzo, *Probes of universal extra dimensions at colliders*, *Phys. Rev.* **D 64** (2001) 095010 [hep-ph/0106336].

[32] T. Stelzer, Z. Sullivan and S. Willenbrock, *Single top quark production at hadron colliders*, *Phys. Rev.* **D 58** (1998) 094021 [hep-ph/9807340].

JHEP03(2002)010

[33] M. Beneke et al., *Top quark physics*, hep-ph/0003033.

[34] F. del Aguila and J. Santiago, *Large top mixing from extra dimensions*, hep-ph/0011143.

[35] ECFA/DESY LC Physics Working Group collaboration, J.A. Aguilar-Saavedra et al., *Tesla technical design report part III. Physics at an e^+e^- linear collider*, hep-ph/0106315.

[36] S. Mouslopoulos, *Bulk fermions in multi-brane worlds*, *J. High Energy Phys.* **05** (2001) 038 [hep-th/0103184].

JHEP03(2002)010