PHYSICAL REVIEW D **85,** 016005 (2012)

Radion flavor violation in warped extra dimensions

K. Huitu,[1] S. Khalil,[2,3] A. Moursy,[2] S. K. Rai,[4] and A. Sabanci[1]

[1]*Department of Physics, and Helsinki Institute of Physics, P.O. Box 64, FIN-00014 University of Helsinki, Finland*

[2]*Centre for Theoretical Physics, The British University in Egypt, El Sherouk City, 11837, Egypt*

[3]*Department of Mathematics, Ain Shams University, Faculty of Science, Cairo, 11566, Egypt*

[4]*Department of Physics, and Oklahoma Center for High Energy Physics, Oklahoma State University Stillwater, Oklahoma 74078, USA*

(Received 1 June 2011; revised manuscript received 2 December 2011; published 12 January 2012)

We analyze the flavor violation in warped extra dimension due to radion mediation. We show that $\Delta S = 2$ and $\Delta B = 2$ flavor violating processes impose stringent constraints on radion mass, m_ϕ and the scale Λ_ϕ. In particular, for $\Lambda_\phi \sim \mathcal{O}(1)$ TeV, $B_d^0 - \bar{B}_d^0$ implies that $m_\phi \gtrsim 65$ GeV. We also study radion contributions to lepton flavor violating processes: $\tau \to (e, \mu)\phi$, $\tau \to e\mu^+\mu^-$ and $B \to l_i l_j$. We show that $\text{BR}(B_s \to \mu^+\mu^-)$ can be of order 10^{-8}, which is reachable at the LHCb. The radion search at LHC, through the flavor violation decays into $\tau\mu$ or top-charm quarks, is also considered.

DOI: 10.1103/PhysRevD.85.016005 PACS numbers: 11.10.Kk, 12.15.Ff, 13.20.−v, 13.25.−k

I. INTRODUCTION

Extra dimensions have been proposed as an alternative way to address the origin of the large scale discrepancy between the Planck scale and the electroweak scale, known as the hierarchy problem [1]. The warped extra dimension is one of the interesting possibilities for a geometrical way to look at this problem. In the Randall-Sundrum (RS) model with two branes [2] the electroweak scale is exponentially suppressed and a large hierarchy between the Planck scale and the TeV scale is obtained. The original RS model is based on the assumption that the standard model (SM) fields are localized to one of the boundaries and gravity only is allowed to propagate in the bulk. In this scenario, the nonrenormalizable operators in the four-dimensional effective theory are only TeV-scale suppressed. This would lead to rapid proton decay and unacceptable flavor violation. If the SM fermions are assumed to be propagating in the bulk one may be able to overcome these problems and also explain the fermion mass hierarchy [3,4]. However, it was shown that in this case flavor changing neutral currents impose strong constraints on the four-dimensional scale [4].

The radius of the extra dimension in the RS model is assumed to be fixed by a given constant. Goldberger and Wise [5] proposed a mechanism to understand the possible mechanism for radius stabilization. It was shown that by adding a scalar field to the bulk, a potential for the radion field is obtained that dynamically generates a vacuum expectation value (VEV) of the radion. This VEV, which is related to the size of the extra dimension, can be naturally of the order of TeV. The radion field $\phi(x)$ arises as the pseudo-Goldstone boson associated with translation symmetry breaking after stabilizing the extra dimension. In this case, the radion mass is given by [5]

$$m_\phi^2 = \frac{k^2 v_v^2}{3M_5^3}\epsilon^2 e^{-2kr_c\pi}, \qquad (1)$$

where M_5 is the five-dimensional Planck scale, $k \simeq M_{\text{Pl}}$, $kr_c \simeq 12$, $\epsilon \ll 1$. Therefore, the mass of radion is typically of the order of a few GeV's. Hence, it may be the lightest new (non-SM) particle in this type of model with warped geometry.

The radion phenomenology has been discussed in several papers [6–9] and recently flavor changing neutral currents mediated by radion field, like $t \to \phi c$ and ϵ_K, have been considered [10]. These analyses showed that the top decay does not impose any significant constraint on the stabilization scale Λ_ϕ, but the *CP* violating parameter ϵ_K may strongly constrain it. However, ϵ_K depends on the *CP* violating phases assumed in the Yukawa matrices, therefore it cannot be used to give a model independent constraint on Λ_ϕ. Our goal in this paper is to pursue this study and consider possible constraints due the experimental bounds of $\Delta S = 2$ and $\Delta B = 2$ processes. In addition, we consider radion contributions to lepton flavor violating processes like $\tau \to (e, \mu)\phi$ and $\tau \to e\mu^+\mu^-$, in addition to $B \to l_i l_j$. We show that although the radion effects enhance the amplitudes of these process, their branching ratios remain below the current experimental limits. We also analyze the search for radion at the LHC. In particular, we focus on the flavor violation decay of radion to $\tau\mu$ or to top-charm quarks.

The paper is organized as follows. In Sec. II, we briefly review the radion interactions with the SM model fermions propagating in the five-dimensional bulk while the Higgs is localized on the TeV brane. We emphasize the radion flavor violating couplings with the SM fermions. Sections III and IV are devoted for analyzing the radion contributions to $\Delta S = 2$ and $\Delta B = 2$ transitions and the constraints imposed on the scale Λ_ϕ and radion mass. It turns out that the $B_d - \bar{B}_d$ mixing gives the strongest bounds on Λ_ϕ and m_ϕ. In Sec. V we study the radion contribution to the decays $B_s \to l_i l_j$. The effects of the radion mediation in lepton flavor violating processes like $\tau \to (e, \mu)\phi$ and $\tau \to e\mu^+\mu^-$ are described in Sec. VI.

The radion search at the LHC is discussed in Sec. VII. Finally, we give our conclusions in Sec. VIII.

II. RADION INTERACTIONS WITH THE SM FERMIONS

We consider the following five-dimensional anti-de Sitter (AdS) space-time [7]:

$$ds^2 = \left(\frac{R}{z}\right)^2 (e^{-2F}\eta_{\mu\nu}dx^\mu dx^\nu - (1+2F)^2 dz^2), \tag{2}$$

where z refers to the conformally flat AdS background with $R' < z < R$. R is the AdS curvature and is given by $R = 1/k \simeq 1/M_{\rm Pl}$ while $R' \simeq 1/\text{TeV}$. The scalar function $F(x,z)$ corresponds to the radion fluctuation around the stabilized radius. From five-dimensional Einstein equations one can show that the metric perturbation $F(x,z)$ is given by

$$F(x,z) = \frac{\phi(x)}{\Lambda_\phi}\left(\frac{z}{R'}\right)^2, \tag{3}$$

where $\Lambda_\phi \equiv \sqrt{6}/R'$. Therefore the square root of the five-dimensional metric determinant $\sqrt{g}$ is given at linear order on F, by

$$\sqrt{g} \approx \left(\frac{R}{z}\right)^5 \left(1 - 2\frac{\phi(x)}{\Lambda_\phi}\left(\frac{z}{R'}\right)^2\right). \tag{4}$$

The five-dimensional action for bulk fermions can be written as

$$S_f = \int d^4x dz \sqrt{g}\Big[\frac{i}{2}(\bar{\mathcal{Q}}_i\Gamma^A\mathcal{D}_A\mathcal{Q}_i - \mathcal{D}_A\bar{\mathcal{Q}}_i\Gamma^A\mathcal{Q}_i) + \frac{c_{q_i}}{R}\bar{\mathcal{Q}}_i\mathcal{Q}_i + (\mathcal{Q}\to\mathcal{U},\mathcal{D}) + (Y^u_{ij}\sqrt{R}\bar{\mathcal{Q}}_i\mathcal{H}\mathcal{U}_j + Y^d_{ij}\sqrt{R}\bar{\mathcal{Q}}_i\mathcal{H}\mathcal{D}_j + \text{h.c.})\Big], \tag{5}$$

where Γ matrices are given by $\Gamma^A = \gamma^a e^A_a$, $a = 0,1,2,3,5$ stands for five-dimensional Lorenz indices. γ^a are the ordinary γ matrices with $\gamma^5 = i\,\text{diag}(1_2, -1_2)$. Here the five-dimensional fermion mass is given in terms of the scale R and the bulk parameter c_f. We assume that the Higgs field is localized on the TeV brane, i.e., $\mathcal{H}(x,z) = H(x)\delta(z-R')$. $\mathcal{Q}_i$, $\mathcal{U}_i$, and $\mathcal{D}_i$ are the five-dimensional fermions, with flavor indices $i, j = 1,2,3$, which contain the four-dimensional SM $SU(2)_L$ doublet and singlet fermions, respectively. They can be written in two component spinor notation as follows:

$$\mathcal{Q}_i = \begin{pmatrix}\mathcal{Q}^i_L\\ \bar{\mathcal{Q}}^i_R\end{pmatrix}, \qquad \mathcal{U}_i = \begin{pmatrix}\mathcal{U}^i_L\\ \bar{\mathcal{U}}^i_R\end{pmatrix}, \quad \text{and} \quad \mathcal{D}_i = \begin{pmatrix}\mathcal{D}^i_L\\ \bar{\mathcal{D}}^i_R\end{pmatrix}. \tag{6}$$

The Kaluza-Klein decomposition for the five-dimensional bulk fields is, as usual, given by

$$\mathcal{Q}_{L,R}(x,z) = \sum_n Q^n_{L,R}(z) q^n_{L,R}(x) \tag{7}$$

with similar expressions for $\mathcal{U}_{L,R}$ and $\mathcal{D}_{L,R}$. The zero modes $q_L(x)$, $u_R(x)$ and $d_R(x)$ define the four-dimensional SM fermions that satisfy the Dirac equations

$$-i\bar{\sigma}^\mu\partial_\mu q^i_L + m^u_{ij}\bar{u}^j_R = 0, \tag{8}$$

$$-i\bar{\sigma}^\mu\partial_\mu q^i_L + m^d_{ij}\bar{d}^j_R = 0, \tag{9}$$

where m^q_{ij} are the mass matrices for up and down quarks which generally are not diagonal in flavor space. Also m^q_{ij} is not simply the induced mass on the TeV brane, given by Higgs VEV and the effective Yukawa coupling weighted by zero mode profiles. m^q_{ij} is the mass eigenvalue that emerges from the solution of the coupled bulk equations of motion, taking into account the Higgs interactions. In general, the physical mass receives corrections from the reaction of the wave functions to the brane where the Higgs is localized [7],

$$(m^q{}_{ij})^2 = (M_D)^2_{ij}\frac{(1-2c^i_L)(1+2c^j_R)}{(1-\lambda^{1-2c^i_L})(1-\lambda^{1+2c^j_R})}, \tag{10}$$

where $\lambda \equiv R/R'$ and M_D is the localized Dirac mass, i.e., induced mass on the brane through the Higgs VEV. Similar expression for charged lepton masses can also be obtained. The boundary conditions are usually chosen such that

$$\mathcal{Q}_R|_{z=R,R'} = \mathcal{U}_L|_{z=R,R'} = \mathcal{D}_L|_{z=R,R'} = 0. \tag{11}$$

These conditions allow the $SU(2)_L$ doublet (left-handed state) and singlet (right-handed state) only to have zero modes. Moreover, due to the arising discontinuities, one should impose the following conditions as well [7]:

$$\mathcal{Q}_R|_{R'^-} = -M^u_D R' \mathcal{U}_R|_{R'^-}, \tag{12}$$

$$\mathcal{U}_L|_{R'^-} = -M^u_D R' \mathcal{Q}_L|_{R'^-}. \tag{13}$$

Similar boundary conditions are applied for $\mathcal{D}_{L,R}$. The bulk wave functions can be found in Ref. [7]. For the zero modes, with the approximation $mR' \ll 1$, the associated wave functions can be written as

$$Q^0_L(z) = z^2 A^q_L\left[\frac{mz^{c^q_L+1}}{2c^q_L+1} + \frac{z^{-c^q_L}R^{2c^q_L-1}}{m}\left(1 - \frac{m^2z^2}{2-4c^q_L}\right)\right],$$
$$Q^0_R(z) = z^2 A^q_R\left[z^{c^q_R}\left(1 - \frac{m^2z^2}{2+4c^q_R}\right) + m^2\frac{R^{1+2c^q_R}}{1+2c^q_R}\frac{z^{1-c^q_R}}{1-2c^q_R}\right]. \tag{14}$$

The parameters $A^q_{L,R}$ can be determined from the normalization conditions:

$$\int dz \left(\frac{R}{z}\right)^4 |Q^0_{L,R}|^2 = 1. \tag{15}$$

Similar expressions can be obtained for $U^0_{L,R}$ and $D^0_{L,R}$. In this respect, the general expression for the Lagrangian of radion interaction with SM fermions is given by [7]

$$\begin{aligned}\mathcal{L}_\phi = &\frac{\phi(x)}{\Lambda_\phi}(q^i_L \bar{u}^j_R + \bar{q}^i_L u^j_R) \times \int dz \left(\frac{R}{z}\right)^2 \frac{R^2}{R'^2} \\ &\times \bigg[-\frac{m^u_{ij}}{2}((Q^i_L)^2 + (Q^i_R)^2) + 2\Big(Q^i_L(Q^i_R)' - (Q^i_L)'Q^j_R \\ &+ \frac{c^{q_i}_L}{z} Q^i_L Q^i_R\Big) + (Q^i_{L,R} \to U^j_{L,R}, D^j_{L,R})\bigg]. \end{aligned} \tag{16}$$

In addition to the approximation $mR' \ll 1$, for light fermions which are usually assumed to be localized near the Planck brane, i.e., $c_L > 1/2$ and $c_R < -1/2$ one finds that the associated radion couplings take the following simple form:

$$\frac{m^{u,d}_{ij}}{\Lambda_\phi}(c^{q_i}_L - c^{u_j,d_j}_R). \tag{17}$$

Transforming to mass eigenstate via the unitary matrices, V^d_L (obtained by left-handed quark rotation) and V^d_R (obtained by right-handed quark rotation) will diagonalize the down mass matrix. In this basis, the radion couplings with down quarks $Y_{\phi d_i d_j}$ are nonuniversal and are given by

$$Y_{\phi d_i d_j} = (V^d_L)^T_{ik} \cdot \frac{m^d_{kl}}{\Lambda_\phi}(c^{q_k}_L - c^{d_l}_R).(V^d_R)_{lj}. \tag{18}$$

It is clear that this flavor violation can be mediated at tree level by the radion propagation, which might be quite dangerous and lead to strong bounds on the stabilization scale Λ_ϕ.

Note that the bulk mass parameters $c^{u,d}_{L,R}$ and five-dimensional Yukawa couplings are free parameters to be fixed by the observable masses and mixing. Therefore, in this class of models the number of free parameters is larger than the number of the quark masses and mixings. In our analysis, as an example, we consider the following values of $c^{u,d}_{L,R}$ that lead to consistent quark masses at the weak scale[1]

$$\begin{aligned} &c^L_{q1} = 0.72, && c^R_{d1} = -0.57, && c^R_{u1} = -0.63 \\ &c^L_{q2} = 0.60, && c^R_{d2} = -0.57, && c^R_{u2} = -0.30 \\ &c^L_{q3} = 0.35, && c^R_{d3} = -0.60, && c^R_{u3} = -0.10 \end{aligned} \tag{19}$$

We also fix the Dirac mass $(M_D)_{ij} = v l_{ij}$, where l_{ij} are dimensionless quantities of order unity obtained from the five-dimensional Yukawa couplings and v is taken to be the SM VEV, namely

$$\begin{aligned} l^d_{ij} &= \begin{pmatrix} 0.50 & -2.00 & -2.00 \\ 1.48 & -0.90 & 2.00 \\ 0.52 & -0.50 & 0.70 \end{pmatrix}, \\ l^u_{ij} &= \begin{pmatrix} 0.80 & -1.90 & -2.00 \\ 1.23 & 1.20 & -1.04 \\ 1.85 & 1.66 & -0.80 \end{pmatrix}. \end{aligned} \tag{20}$$

In general, these parameters are complex. However, the corresponding phases may lead to a large contribution to the CP violating processes (as ϵ_K), which is inconsistent with the SM expectations. Therefore, these phases are typically constrained and set to zero unless one assumes a specific texture of flavor that suppresses both CP conserving and CP violating flavor changing effects as in Ref. [11].

Using these parameters, one obtains the following quark masses:

$$\begin{aligned} &m_d = 2.7\ \text{MeV}, && m_u = 1.66\ \text{MeV}, \\ &m_s = 47\ \text{MeV}, && m_c = 1.2\ \text{GeV}, \\ &m_b = 3.524\ \text{GeV}, && m_t = 171.25\ \text{GeV} \end{aligned} \tag{21}$$

and the Cabibbo-Kobayashi-Maskawa (CKM) matrix is given by

$$V_{\text{CKM}} = \begin{pmatrix} 0.972\,061 & 0.234\,68 & 0.004\,733\,66 \\ 0.234\,685 & 0.971\,301 & 0.038\,695\,6 \\ 0.004\,483\,28 & 0.038\,725\,4 & 0.999\,24 \end{pmatrix}. \tag{22}$$

Corresponding equations can be written for leptons as well, and thus acceptable lepton masses can be derived from Eq. (10) where q has been replaced by l, using the following parameters

$$\begin{aligned} &c^L_{l1} = 0.705, \quad c^e_R = -0.52 \quad c^L_{l2} = 0.655, \\ &c^\mu_R = -0.53 \quad c^L_{l3} = 0.550, \quad c^\tau_R = -0.585, \\ &l^l_{ij} = \begin{pmatrix} -1.37 & 1.19 & 1.04 \\ 1.10 & 1.10 & 0.90 \\ -0.80 & 0.90 & 0.90 \end{pmatrix}. \end{aligned} \tag{23}$$

In this case, one finds $m_e = 0.511$ MeV, $m_\mu = 0.106$ GeV, and $m_\tau = 1.777$ GeV. Also the Yukawa couplings of radion-fermion-antifermion can be approximately written, in terms of the scale Λ_ϕ, as

[1]We modify the model in [4] by imposing the conditions for c parameters of light fermions, i.e., $c_L > 1/2$ and $c_R < -1/2$.

$$Y_u = \frac{1\ \text{GeV}}{\Lambda_\phi}\begin{pmatrix} 0.002\,242\,09 & 0.020\,937\,9 & 0.146\,558 \\ -0.003\,113\,23 & 0.919\,45 & 0.002\,242\,09 \\ -0.445\,271 & 15.937 & 100.47 \end{pmatrix}, \tag{24}$$

$$Y_d = \frac{1\ \text{GeV}}{\Lambda_\phi}\begin{pmatrix} 0.003\,473\,54 & -0.000\,530\,211 & -0.006\,184\,97 \\ -0.000\,427\,834 & 0.055\,131 & -0.054\,655\,9 \\ 0.034\,730\,5 & -0.013\,863\,9 & 3.260\,96 \end{pmatrix}, \tag{25}$$

and

$$Y_e = \frac{1\ \text{GeV}}{\Lambda_\phi}\begin{pmatrix} 0.000\,657\,813 & -0.000\,167\,775 & -0.003\,912\,32 \\ -0.000\,800\,194 & 0.124\,867 & 0.001\,014\,46 \\ 0.015\,131\,8 & -0.010\,937\,1 & 1.912\,15 \end{pmatrix}. \tag{26}$$

The c and l parameters found here are obviously not unique and one may wonder how general our results are using these sets. To study that, we have generated another parameter set both for quarks and leptons. Although the c and l values in the new sets are clearly different from the ones shown here and used in the analyses of the latter sections, it turns out that the results remain qualitatively the same, and quantitatively change only little.

III. RADION CONTRIBUTION TO $\Delta S = 2$ TRANSITIONS

We start our analysis for radion flavor violation by considering the radion contribution to $\Delta S = 2$ processes, where S refers to the s-quark number, in particular, to $K^0 - \bar{K}^0$. Generically, the $K_L - K_S$ mass difference ΔM_K is defined as

$$\Delta M_K = 2|\langle K|H_{\text{eff}}^{\Delta S=2}|\bar{K}\rangle|, \tag{27}$$

where $H_{\text{eff}}^{\Delta S=2}$ is the effective Hamiltonian for $\Delta S = 2$ transition. With radion contribution to the off-diagonal entry in the K meson, the mass matrix $\mathcal{M}_{12}(K) = \langle K|H_{\text{eff}}^{\Delta S=2}|\bar{K}\rangle$ is given by

$$\mathcal{M}_{12}(K) = \mathcal{M}_{12}^{\text{SM}}(K) + \mathcal{M}_{12}^{\text{rad}}(K). \tag{28}$$

Here $\mathcal{M}_{12}^{\text{SM}}(K)$ is the SM contribution and is given by

$$\mathcal{M}_{12}^{\text{SM}}(K) = \frac{G_F^2}{12\pi^2}\hat{B}_K f_K^2 M_K M_W^2(\eta_1(\lambda_c^*)^2 S_0(x_c) + \eta_2(\lambda_t^*)^2 S_0(x_t) + 2\eta_2(\lambda_t^*)(\lambda_c^*)S_0(x_c, x_t)) \tag{29}$$

where $\lambda_i = V_{is}^* V_{id}$ and other parameters and loop functions which appear in the above equation can be found in Ref. [12]. The SM expectation for ΔM_K is given by

$$\Delta M_K^{\text{SM}} = 2.7018 \times 10^{-15}\ \text{GeV}, \tag{30}$$

which lies in the ballpark of the measured value [13]:

$$\Delta M_K^{\text{exp}} = 3.483 \pm 0.006 \times 10^{-15}\ \text{GeV}. \tag{31}$$

However, a precise prediction cannot be made due to the hadronic and CKM uncertainties.

Unlike the SM, the radion contribution to the $K^0 - \bar{K}^0$ mixing is at the tree level, as shown in Fig. 1. The corresponding effective Hamiltonian is given by

$$H_{\text{eff}}^{\text{rad}}(\Delta S = 2) = \sum_{i=1,2}(C_i Q_i + \tilde{C}_i \tilde{Q}_i), \tag{32}$$

where C_i, $\tilde{C}_i$, Q_i and $\tilde{Q}_i$ are the Wilson coefficients and operators with

$$\begin{aligned} Q_1 &= (\bar{s}_L d_R)(\bar{s}_L d_R), \\ Q_2 &= (\bar{s}_L d_R)(\bar{s}_R d_L), \\ C_1 &= -\frac{Y_{\phi s_L d_R}^2}{M_K^2 - m_\phi^2 + i m_\phi \Gamma_\phi}, \\ C_2 &= -\frac{Y_{\phi s_L d_R} Y_{\phi s_R d_L}}{M_K^2 - m_\phi^2 + i m_\phi \Gamma_\phi}. \end{aligned} \tag{33}$$

The operators $\tilde{Q}_i$ and Wilson coefficients $\tilde{C}_i$ are obtained from Q_i and C_i by exchanging $L \leftrightarrow R$. Note that for $m_\phi \gg M_K$, the Wilson coefficients are given by $\sim Y^2/M_\phi^2$. For $m_\phi < M_K$, if we assume that the momentum transfer in the four-fermion operator is around M_K the coefficients are $\sim Y^2/M_K^2$, which is a consistent approximation since for light radion the $K - \bar{K}$ transition occurs through the decay of K into ϕ and X_d.

The mass of the radion is in the range of a few GeVs when the external momenta are neglected. The matrix

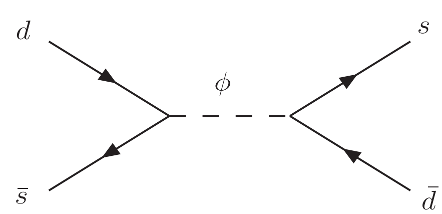

FIG. 1. Radion contributions to the $K^0 - \bar{K}^0$ mixing.

elements of the operators Q_i between K mesons in the vacuum insertion approximation (VIA) are given by [14]:

$$\langle \bar{K}^0|Q_1|K^0\rangle_{\rm VIA} = -\frac{5}{24}\left(\frac{M_K}{m_s+m_d}\right)^2 M_K f_K^2,$$
$$\langle \bar{K}^0|Q_2|K^0\rangle_{\rm VIA} = \left[\frac{1}{24}+\frac{1}{4}\left(\frac{M_K}{m_s+m_d}\right)^2\right] M_K f_K^2, \qquad (34)$$

where m_s and m_d are the masses of s and d quarks, respectively. In the case of the renormalized operators, we define the B parameters as

$$\langle \bar{K}^0|\hat{Q}_1(\mu)|K^0\rangle = -\frac{5}{24}\left(\frac{M_K}{m_s(\mu)+m_d(\mu)}\right)^2 M_K f_K^2 B_1(\mu),$$
$$\langle \bar{K}^0|\hat{Q}_2(\mu)|K^0\rangle = \frac{1}{4}\left(\frac{M_K}{m_s(\mu)+m_d(\mu)}\right)^2 M_K f_K^2 B_2(\mu), \qquad (35)$$

where $\hat{Q}_i(\mu)$ denotes the operators renormalized at the scale μ. For the scale $\mu = 2$ GeV, $B_1(\mu)$ and $B_2(\mu)$ are given by [14]:

$$B_1(\mu) = 0.66(4), \qquad B_2(\mu) = 1.03(6). \qquad (36)$$

Using Eq. (18) and the values of the c parameters in Eq. (19), we can compute the values of the Wilson coefficients at the scale of the radion mass. Since the K decay occurs at 2 GeV, we should run the Wilson coefficients from the scale of the radion mass to the scale of 2 GeV, considering all thresholds, using the following general renormalization group equations [15] that runs the Wilson coefficients from scale M to another scale μ

$$C_i(\mu) = \left[1+\frac{\alpha_s(\mu)-\alpha_s(M)}{4\pi}J_f\right]\left[\frac{\alpha_s(M)}{\alpha_s(\mu)}\right]^{d_f} C_i(M), \qquad (37)$$

where

$$d_f = \frac{\gamma^{(0)}}{2\beta_0}, \qquad J_f = \frac{d_f}{\beta_0}\beta_1 - \frac{\gamma^{(1)}}{2\beta_0},$$

and

$$\gamma^{(0)} = 6\frac{N-1}{N},$$
$$\gamma^{(1)} = \frac{N-1}{2N}\left[-21+\frac{57}{N}-\frac{19}{3}N+\frac{4}{3}f\right],$$
$$\beta_0 = (11N-2f)/3, \qquad (38)$$
$$\beta_1 = \frac{34}{3}N^2 - \frac{10}{3}Nf - \frac{N^2-1}{N}f,$$

where N is the number of colors and f is the number of active flavors. Also we run the produced masses, Eq. (21), from the weak scale to the scale of 2 GeV.

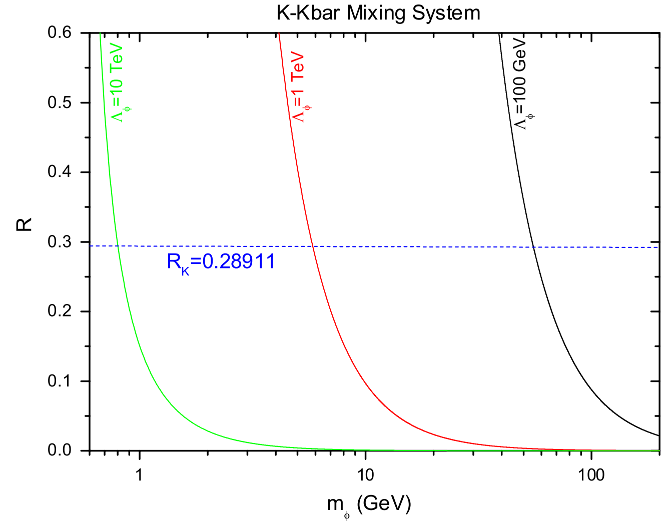

FIG. 2 (color online). The ratio $R_K = |\frac{\mathcal{M}_{12}^{\rm rad}(K)}{\mathcal{M}_{12}^{\rm SM}(K)}|$ as function of radion mass m_ϕ for scale $\Lambda_\phi = 0.1,\ 1,\ 10$ TeV.

Using the above expressions, one can compute the radion contribution to ΔM_K. The experimental limits of ΔM_K implies that $R_K = |\frac{\mathcal{M}_{12}^{\rm rad}(K)}{\mathcal{M}_{12}^{\rm SM}(K)}| \leq 0.289\,11$, which leads to an upper bound on the radion contribution. This upper bound imposes a stringent constraint on the scale Λ_ϕ and the radion mass m_ϕ. In Fig. 2 we show the constraint on the radion mass m_ϕ, due to the $K-\bar{K}$ mixing system, for three values of the scale Λ_ϕ: 0.1, 1, and 10 TeV. As can be seen from this figure, a very light radion $\mathcal{O}(1)$ GeV can be allowed if Λ_ϕ of order 10 TeV. However $\Lambda_\phi \lesssim 1$ TeV can be consistent with ΔM_K experiment bound if $m_\phi \geq \mathcal{O}(5)$ GeV.

In order to study the sensitivity of these bounds on m_ϕ and Λ_ϕ to the values of the bulk mass parameters c and five-dimensional Yukawa parameters l, we consider another example of these parameters that produce the correct quark masses and $V_{\rm CKM}$ mixing matrix. Namely, the following set of parameters is considered:

$$c_{q1}^L = 0.643, \quad c_{d1}^R = -0.675, \quad c_{u1}^R = -0.645$$
$$c_{q2}^L = 0.583, \quad c_{d2}^R = -0.630, \quad c_{u2}^R = -0.630$$
$$c_{q3}^L = 0.317, \quad c_{d3}^R = -0.590, \quad c_{u3}^R = 0.150 \qquad (39)$$

$$l_{ij}^d = \begin{pmatrix} 1.1 & 0.49 & 0.88 \\ 0.86 & -0.96 & -0.53 \\ 0.99 & 1.1 & -1.20 \end{pmatrix},$$
$$l_{ij}^u = \begin{pmatrix} -0.44 & 1.21 & -0.50 \\ -0.91 & -0.24 & 1.22 \\ 0.40 & -1.15 & 0.99 \end{pmatrix}. \qquad (40)$$

Although these parameters are different from the ones listed in Eq. (19) and (20), we find out that there is no important difference between the two examples and the above bounds derived on m_ϕ and Λ_ϕ remains intact.

IV. RADION CONTRIBUTION TO $\Delta B = 2$ TRANSITIONS

There are two neutral $B^0 - \bar{B}^0$ meson systems: $B_q^0 - \bar{B}_q^0$, with $q = d, s$. In this systems, the flavor eigenstates are given by $B_q = (\bar{b}q)$ and $\bar{B}_q = (b\bar{q})$. As in the $K^0 - \bar{K}^0$ system, the mass difference between mass eigenstates B_L^q and B_H^q is defined as

$$\Delta M_{B_q} = M_{B_H}^q - M_{B_L}^q = 2|\mathcal{M}_{12}^q| = 2|\langle B_q^0|H_{\rm eff}^{\Delta B=2}|\bar{B}_q^0\rangle|. \tag{41}$$

The experimental values for mass difference for $\Delta M_{B_d}^{\rm exp}$ and $\Delta M_{B_s}^{\rm exp}$ are given by [13]

$$\Delta M_{B_d}^{\rm exp} = (3.337 \pm 0.033) \times 10^{-13}\ {\rm GeV}, \tag{42}$$

$$\Delta M_{B_s}^{\rm exp} = (117.0 \pm 0.8) \times 10^{-13}\ {\rm GeV}. \tag{43}$$

The SM contribution for ΔM_{B_q} at next-to-leading order is given by [16]

$$\Delta M_{B_q}^{\rm SM} = \frac{G_F^2}{6\pi^2}\eta_B m_{B_q}\hat{B}_{B_q}F_{B_q}^2 M_W^2 S_0(x_t)(V_{tq}^* V_{tb})^2, \tag{44}$$

where F_{B_q} is the B_q meson decay constant for $q = d, s$ and $\hat{B}_{B_q}$ is the renormalization-group invariant parameters [17]. One can show that the SM predictions for ΔM_{B_q} are given by

$$\begin{aligned}\Delta M_{B_d}^{\rm SM} &= 3.581\,87 \times 10^{-13}\ {\rm GeV},\\ \Delta M_{B_s}^{\rm SM} &= 104.19 \times 10^{-13}\ {\rm GeV}.\end{aligned} \tag{45}$$

The leading diagrams of radion contributions to the effective Hamiltonian approach $H_{\rm eff}^{\Delta B_q=2}$ are given by tree level diagrams similar to the diagram of $K^0 - \bar{K}^0$ mixing, with replacing s-quark by b-quark and d-quark by q-quark. The induced effective Hamiltonian for $\Delta B = 2$ radion mediated process is given by

$$H_{\rm eff}^{\rm rad}(\Delta B_q = 2) = \sum_{i=1,2}(C_i Q_i + \tilde{C}_i\tilde{Q}_i), \tag{46}$$

where the operators Q_i and the Wilson coefficients C_i are given by

$$\begin{aligned}Q_1 &= (\bar{b}_L q_R)(\bar{b}_L q_R),\\ Q_2 &= (\bar{b}_L q_R)(\bar{b}_R q_L),\\ C_1 &= \frac{Y_{\phi b_L q_R}^2}{M_{B_q}^2 - m_\phi^2 + i m_\phi \Gamma_\phi},\\ C_2 &= \frac{Y_{\phi b_L q_R} Y_{\phi b_R q_L}}{M_{B_q}^2 - m_\phi^2 + i m_\phi \Gamma_\phi}.\end{aligned} \tag{47}$$

The operators $\tilde{Q}_i$ and the coefficients $\tilde{C}_i$ are obtained from Q_i and C_i by exchanging $L \leftrightarrow R$. Here, all the approximations on Wilson coefficients in Sec. III can be applied via replacing M_K by M_{B_q}. Thus, the renormalized hadronic matrix elements for radion mediated process can be found as [18]:

$$\begin{aligned}\langle\bar{B}_q|\hat{Q}_1(\mu)|B_q\rangle &= -\frac{5}{24}\left(\frac{m_{B_q}}{m_b(\mu) + m_q(\mu)}\right)^2 m_{B_q}F_{B_q}^2 B_1(\mu),\\ \langle\bar{B}_q|\hat{Q}_2(\mu)|B_q\rangle &= \frac{1}{4}\left(\frac{m_{B_q}}{m_b(\mu) + m_q(\mu)}\right)^2 m_{B_q}F_{B_q}^2 B_2(\mu).\end{aligned} \tag{48}$$

Here we adopt the numerical values of B_1, B_2, m_{B_q}, and F_{B_q} as in [18]. Also after calculating the Wilson coefficients at the scale of the radion mass, we derive the corresponding coefficients at $\mu = 4.6$ GeV via Eq. (37). We consider the allowed upper bounds on $\Delta M_{B_q}^{\rm exp}$ to derive new constraints on the radion free parameters in our analysis: Λ_ϕ and m_ϕ.

In Fig. 3 we present the constraints imposed on m_ϕ from the experimental results (using central values of the results) for $B_d^0 - \bar{B}_d^0$ and $B_s^0 - \bar{B}_s^0$ mixing systems, for $\Lambda_\phi = 0.1$, 1, 10 TeV. This figure shows that the experimental limits of $B_d^0 - \bar{B}_d^0$ gives more stringent constraints on $\Lambda_\phi - m_\phi$ than the limits of $B_s^0 - \bar{B}_s^0$ and $K^0 - \bar{K}^0$. In this respect, it is clear that the processes of $\Delta B = 2$ and $\Delta S = 2$ flavor violation play important role in constraining the radion mass and it is no longer a free parameter. For instance if we require that $\Lambda_\phi \sim \mathcal{O}(1)$ TeV, which is favored by solving the hierarchy problem, one finds that the radion mass has the following lower bound: $m_\phi \gtrsim 65$ GeV. We have checked that these bounds are not sensitive to the values of the bulk mass c parameters and five-dimensional l-Yukawa parameters. We obtained very close limits on Λ_ϕ and m_ϕ when we considered the example in Eqs. (39) and (40).

To our knowledge, it is the first time that such a lower bound on radion mass is derived. Nevertheless, if a larger value of the scale Λ_ϕ is considered, i.e., $\Lambda_\phi \gtrsim 10$ TeV, a smaller radion mass, $m_\phi \sim 10$ GeV, can be allowed. As we will show below, a very light radion scenario is stringently constrained by the lepton flavor violation decays $\ell_i \to \ell_j\phi$.

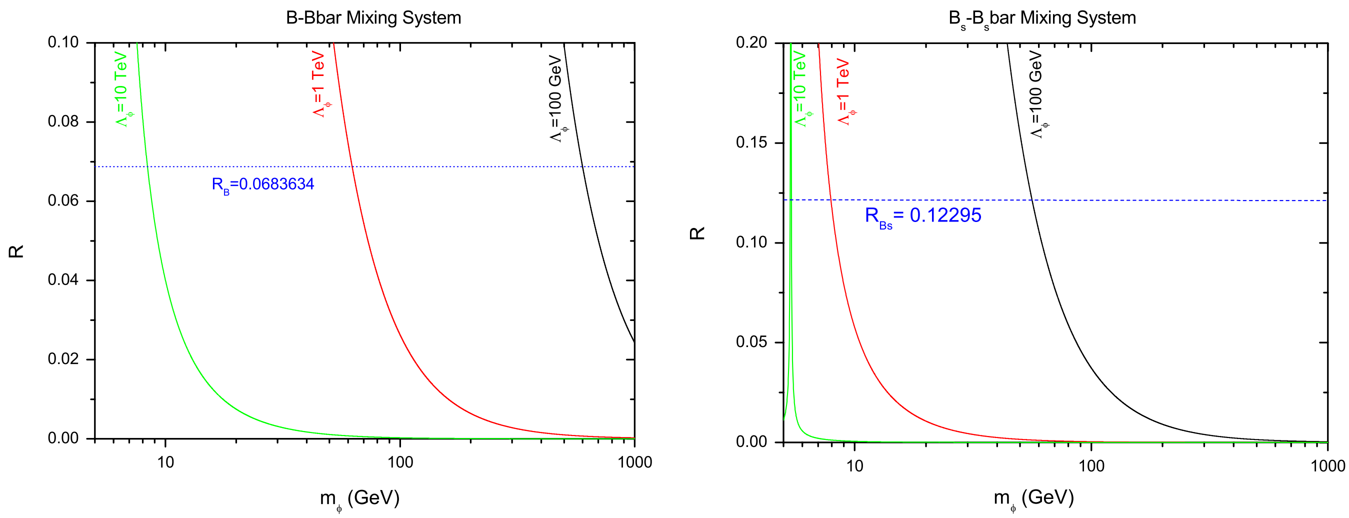

FIG. 3 (color online). The constraints imposed on the radion mass m_ϕ, due to the radion contribution to $B_d^0 - \bar{B}_d^0$ and $B_s^0 - \bar{B}_s^0$ mixing systems, for $\Lambda_\phi = 0.1$, 1, 10 TeV.

V. RADION CONTRIBUTION TO LEPTONIC *B*-DECAYS

We now consider the radion contribution to the leptonic B-decays: $B_q \to l_i l_j$, where $q \equiv d, s$ [Fig. 4]. In this class of models with warped geometry, the B-meson decay into leptons can be generated at tree level through radion exchange. Note that, as in the quark sector, the radion couplings with the charged leptons are given by

$$Y_{\phi l_i l_j} = \frac{m^l_{ij}}{\Lambda_\phi}(c_L^{l_i} - c_R^{e_j}), \tag{49}$$

which transforms to the following expression in lepton mass basis:

$$Y_{\phi l_i l_j} = (V_L^l)^T_{ik} \frac{m^l_{kk'}}{\Lambda_\phi}(c_L^{l_k} - c_R^{e_{k'}})(V_R^l)_{k'j}. \tag{50}$$

It is worth noting that these nonuniversal couplings are obtained due to the mismatch between the diagonalization of charged lepton mass matrix m^l_{ij} and the charged lepton-radion couplings $\sim m^l_{ij}(c_L^{l_i} - c_R^{e_j})$. The transition amplitude of this process is given by

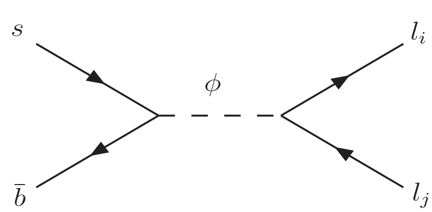

FIG. 4. Radion contributions to the $B_s \to l_i l_j$.

$$\mathcal{A}(\bar{B}_q \to l_i^- l_j^+) = \left[\frac{1}{q^2 - m_\phi^2}\right][Y_1 Y_3 \bar{l}_i P_R l_j \langle 0|\bar{s} P_L b|B_q\rangle + Y_1 Y_4 \bar{l}_i P_R l_j \langle 0|\bar{s} P_R b|B_q\rangle + Y_2 Y_3 \bar{l}_i P_L l_j \langle 0|\bar{s} P_L b|B_q\rangle + Y_2 Y_4 \bar{l}_i P_L l_j \langle 0|\bar{s} P_R b|B_q\rangle], \tag{51}$$

where the Yukawa couplings Y_i, $i = 1, \ldots, 4$ are defined as

$$Y_1 \equiv Y_{\phi l_{iL}^- l_{jR}^+}, \qquad Y_2 \equiv Y_{\phi l_{iR}^- l_{jL}^+}, \tag{52}$$

$$Y_3 \equiv Y_{\phi b_L q_R}, \qquad Y_4 \equiv Y_{\phi b_R q_L}. \tag{53}$$

The hadronic matrix elements are characterized by the decay constant of the pseudoscalar meson B_q and can be written as [19]

$$\langle 0|\bar{s}\gamma_5 b|B_s(p)\rangle = -i f_{B_s}\frac{M_{B_s}^2}{m_b + m_s}. \tag{54}$$

The partial decay width for the leptonic decay of B_q meson $\Gamma(\bar{B}_q \to \ell_i \ell_j)$ is given by

$$\Gamma(\bar{B}_q \to \ell_i \ell_j) = \frac{1}{16\pi m_{B_q}^3}\overline{|\mathcal{A}|^2}\lambda^{1/2}(m_{\ell_i}^2, m_{\ell_j}^2, m_{B_q}^2), \tag{55}$$

where $\overline{|\mathcal{A}|^2}$ is the spin-averaged amplitude for the radion contribution to the decay and

$$\lambda(x, y, z) = x^2 + y^2 + z^2 - 2xy - 2yz - 2zx.$$

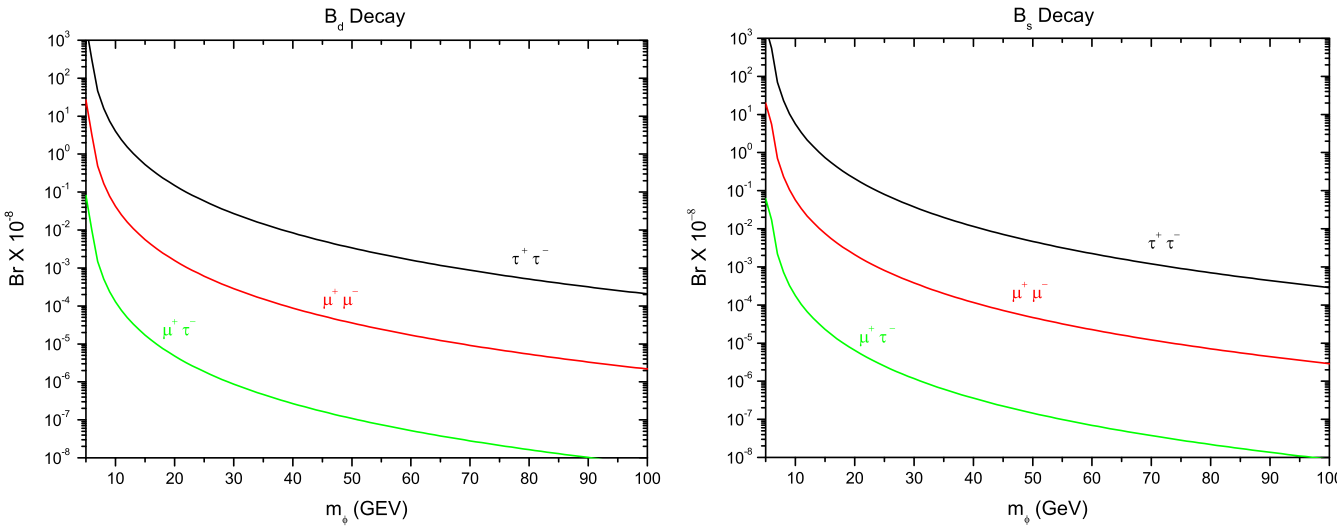

FIG. 5 (color online). The branching ratios of the decays of B_s and B_d to $\mu^+\mu^-$, $\tau^+\tau^-$ and $\mu^\pm\tau^\mp$.

The experimental limits on the branching ratios are given as in the table below.

Br	Experimental limit
$B_s \to \mu^+\mu^-$	$<4.7 \times 10^{-8}$
$B_d \to \mu^+\mu^-$	$<1.5 \times 10^{-8}$
$B_d \to \tau^+\tau^-$	$<4.1 \times 10^{-3}$
$B_d \to \mu^+\tau^-$	$<2.2 \times 10^{-5}$

In Fig. 5 we display the branching ratio of B_q decays to $\mu^+\mu^-$, $\tau^+\tau^-$ and $\mu^\pm\tau^\mp$ as a function of m_ϕ for $\Lambda_\phi = 1$ TeV. From this figure, one finds that the $\mathrm{Br}(B_s \to \mu^+\mu^-)$ can be of order 10^{-8}, i.e., within the range of accessibility at the LHCb, if the radion mass is less than 10 GeV. Also the present LHCb experimental limit: $\mathrm{Br}(B_s \to \mu^+\mu^-) < 6.5 \times 10^{-8}(95\%\ \mathrm{C.L.})$ implies that $m_\phi > 6$ GeV. In addition, it is predicted that within the region of light radion mass the $\mathrm{Br}(B_s \to \tau^+\tau^-)$ is of order 10^{-6}, which can be probed at the LHCb.

VI. RADION CONTRIBUTION TO LEPTON FLAVOR VIOLATING LEPTON DECAYS

In this section we study lepton flavor violating processes in which radion is either a decay product, as in $\ell_i \to \ell_j\phi$, or which is mediated by a radion, e.g. $\ell_i \to \ell_j\ell_k\ell_l$.

A. τ decay to a lepton and radion

We start by studying lepton flavor violation via the process $\tau \to (e, \mu)\phi$. We do not specify the decay products of the radion. It dominantly decays to a gluon pair, but can decay also to a muon pair or a kaon pair, and with a small probability also to an electron pair or a pion pair. Thus we have a muon or an electron from a tau decay, with no missing energy. The limitation of this process is that we can only study radions which are lighter than τ.

The amplitude for the decay is given by

$$\mathcal{A}(\ell_i(k) \to \ell_j(p)\phi(q) = Y_{\phi\bar{\ell}_{Rj}\ell_{Li}}\bar{\ell}_j P_L \ell_i + Y_{\phi\bar{\ell}_{Lj}\ell_{Ri}}\bar{\ell}_j P_R \ell_i. \tag{56}$$

The total partial decay width for the lepton flavor violating decay is

$$\Gamma(\ell_i \to \ell_j\phi)) = \frac{1}{16\pi m_{\ell_i}^3}\overline{|\mathcal{A}|^2}\lambda^{1/2}(m_{\ell_i}^2, m_{\ell_j}^2, m_\phi^2), \tag{57}$$

where M_ϕ is the radion mass and λ is defined after the Eq. (55). The spin-averaged amplitude in the rest frame of the decaying lepton is

$$\overline{|\mathcal{A}|^2} = (Y^2_{\phi\bar{\ell}_{Lj}\ell_{Ri}} + Y^2_{\phi\bar{\ell}_{Rj}\ell_{Li}})\frac{1}{2}(m_{\ell_i}^2 + m_{\ell_j}^2 - M_\phi^2) + 2\mathrm{Re}(Y^*_{\phi\bar{\ell}_{Lj}\ell_{Ri}} Y_{\phi\bar{\ell}_{Rj}\ell_{Li}})m_{\ell_i}m_{\ell_j}. \tag{58}$$

The experimental value for $\Gamma(\tau \to (e, \mu)\phi)$, which can be calculated from the tau life time [13], is found as

$$\Gamma_\tau \simeq (2.259\,692 \pm 0.007\,77) \times 10^{-12}\ \mathrm{GeV}. \tag{59}$$

We have calculated the radion masses and Λ_ϕ which are allowed by the experimental error bars. The result is plotted in Fig. 6. We see that if the scale Λ_ϕ is large, this decay mode may still constrain the radion masses. For example, at 2σ level if $\Lambda_\phi = 10$ TeV, we found that the radion mass must be larger than 1.3 GeV. Thus, if we take $m_\phi \sim 1$ GeV we obtain $\Lambda_\phi \sim 15$ TeV. We have checked that this value respects the condition Eq. (12) in Ref. [20].

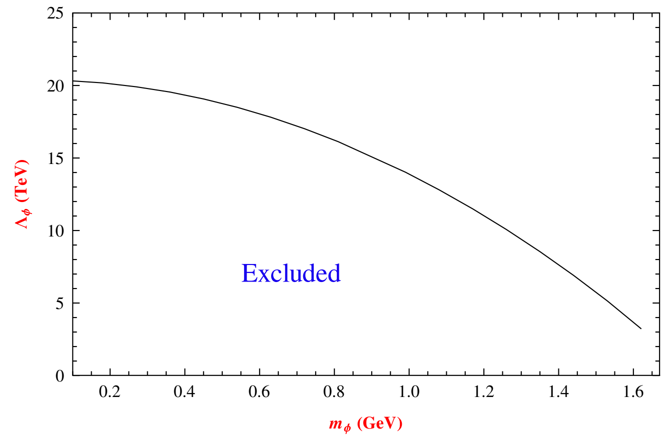

FIG. 6 (color online). The excluded region in (m_ϕ, Λ_ϕ) plane from the $\tau \to (e, \mu)\phi$ decay at 2σ level.

B. τ decay to three leptons

Much studied lepton flavor violating precision measurements include $\ell_i \to \ell_j \ell_k \ell_l$ and $\ell_i \to \ell_j \gamma$. As an example of this class of processes we study $\tau \to e\mu\mu$, shown in Fig. 7, which has the largest coupling constants in our parameter set for leptons. The transition amplitude for the process $\tau \to e\phi^* \to e\mu\mu$ is given by

$$\begin{aligned}\mathcal{A}(\tau(p) \to e^-(p_1)\mu^-(p_2)\mu^+(p_3)) \\ = \frac{1}{q^2 - m_\phi^2 + im_\phi\Gamma_\phi} \sum_{i,j,k,l=L,R i\neq j,k\neq l} \\ \times [Y_{\phi\mu_i\mu_j^+} Y_{\phi\tau_k e_l} (\bar{u}(p_2)P_j v(p_3))(\bar{u}(p_1)P_k u(p))]. \quad (60)\end{aligned}$$

Thus, one can show that the decay rate is given by

$$\begin{aligned}\Gamma(\tau \to e\phi^* \to e\mu\mu) \\ = \frac{m_\tau^5}{3 \times 2^{12}\pi^3} \left(\frac{1}{m_\tau^2 - m_\phi^2 + im_\phi\Gamma_\phi}\right)^2 \\ \times [Y^2_{\phi\mu_L\mu_R^+} Y^2_{\phi\tau_R e_L} + Y^2_{\phi\mu_L\mu_R^+} Y^2_{\phi\tau_L e_R} + Y^2_{\phi\mu_R\mu_L^+} Y^2_{\phi\tau_R e_L} \\ + Y^2_{\phi\mu_R\mu_L^+} Y^2_{\phi\tau_L e_R}]. \quad (61)\end{aligned}$$

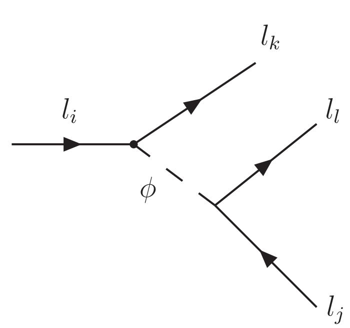

FIG. 7. The decay $\ell_i \to \ell_j \ell_k \ell_l$ through the exchange of radion.

The experimental limit for the branching ratio is [13]:

$$\mathrm{Br}(\tau \to e\mu\mu) < 3.7 \times 10^{-8}. \quad (62)$$

We show the branching ratio as a function of radion masses for $\Lambda_\phi = 0.1$ and 1 TeV in Fig. 8. It is seen that this decay cannot compete with the other discussed ones in restricting Λ_ϕ vs m_ϕ. With our parameter set, the same conclusion seems to hold for other $\ell_i \to$ three leptons and also for lepton gamma modes.

VII. RADION SEARCH AT THE LHC

The radion coupling to the fermions is proportional to the fermion mass. The dominant production mode for the radion at LHC is through gluon fusion. The radion has an enhanced coupling with gluons through the trace anomaly [8]:

$$\mathcal{L}_{\mathrm{int}} = \frac{\phi(x)}{\Lambda_\phi} T^\mu_\mu, \quad (63)$$

with T^μ_μ defined as

$$T^\mu_\mu = \frac{\beta_{\mathrm{QCD}}}{2g_s} \mathrm{Tr}(F^a_{\mu\nu} F^{a\mu\nu}), \quad (64)$$

where $F^a_{\mu\nu}$ is the field strength tensor of $SU(3)$ interactions and β_{QCD} is the QCD beta-function coefficient defined as

$$\frac{\beta_{\mathrm{QCD}}}{2g_s} = -\frac{\alpha_s}{8\pi} b_{\mathrm{QCD}}. \quad (65)$$

Additional contributions coming from the heavy quark loop diagram is suppressed and neglecting that, the cross section of radion production through gluon fusion at LHC can be written as

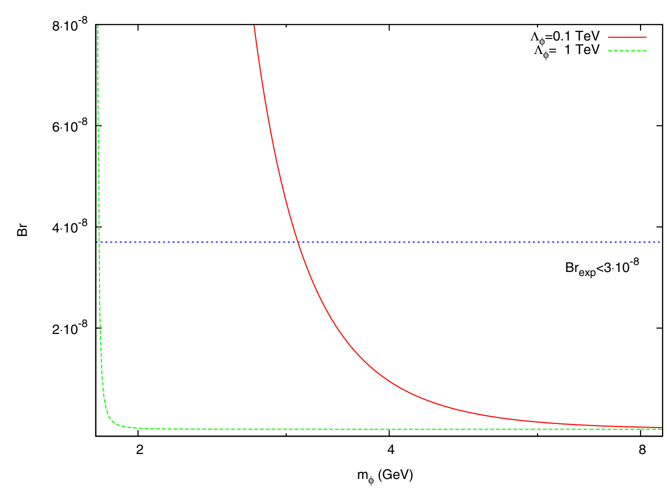

FIG. 8 (color online). Branching ratio of $\tau \to e\mu\mu$ as a function of m_ϕ for the scale $\Lambda = 0.1$ and 1 TeV.

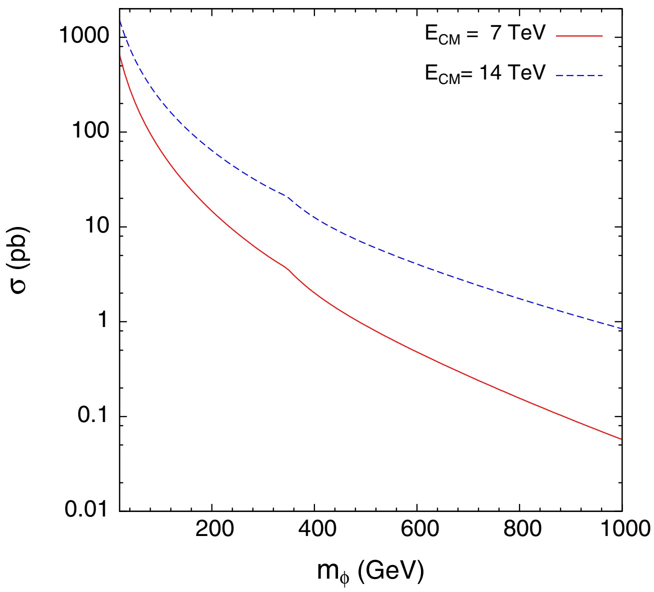

FIG. 9 (color online). Radion production cross section at the LHC for two different center-of-mass energies as a function of the radion mass. We have fixed $\Lambda_\phi = 1$ TeV.

$$\sigma_{LO}(gg \to \phi) = \int_\tau^1 \frac{dx}{x} G(x) G\left(\frac{\tau}{x}\right) \frac{\tau \alpha_s^2}{256\pi\Lambda_\phi^2} |b_{\rm QCD}|^2 \tag{66}$$

where $b_{\rm QCD} = 11 - 2n_f/3$, n_f being the number of quark flavors. The gluon flux in the parton density functions (PDF) is given by $G(x)$ where x is the momentum fraction carried by the gluons. For a radion of mass m_ϕ and the center-of-mass energy $\sqrt{s}$ we define $\tau = m_\phi^2/s$. We calculate the total leading order cross section for the radion production at LHC for two different center-of mass energies, 7 TeV and 14 TeV as a function of the mass of the radion (m_ϕ) with $\Lambda_\phi =$ 1 TeV and is shown in Fig. 9. We use the CTEQ6l PDF [21] for our calculation and the QCD scale Q is set as the radion mass. Note that the cross section scales as $1/\Lambda_\phi^2$. Thus if we reduce Λ_ϕ by a factor 2 then the cross section is increased by a factor of 4. We have already shown that flavor physics constrains the parameter space with lower bounds obtained on the radion mass for $\Lambda_\phi \sim \mathcal{O}(1)$ TeV. For a 100 GeV radion, Λ_ϕ can be as low as 300 GeV which implies a cross section of ~710 pb for the radion production at LHC with $\sqrt{s} = 7$ TeV. For radion of mass less than 100 GeV, there are additional constraints on Λ_ϕ from LEP data [6].

After its production, radion will decay either into gg, W^+W^-, ZZ, $q_i\bar{q}_j$ or $\ell_i^+\ell^-$. Although the first three channels dominate the radion decay [8], one can try to search for flavor violating radion decays through the leptonic decay modes. The large cross section for the radion production can give significant events for the alternative decay channels ($\ell_i^+\ell^-$) which will have smaller SM background. Therefore, they could be striking signatures for radion search at the LHC. However, the leptonic branchings of the radion are very suppressed and fall rapidly with increasing radion mass. The flavor violating leptonic decay channels ($\mu\tau/e\tau$) are further suppressed with branching probabilities even smaller than the diphoton channel and is of the order of 10^{-6} for light radion of mass less than 100 GeV. The $\tau^+\tau^-$ decay channel is about 5% for a 50 GeV radion and $\Lambda_\phi = 500$ GeV. With a good τ-id at LHC, this can be an important channel for the light radion signal at LHC. For a heavier radion with mass greater than the top mass, the radion can decay to a top-quark and charm quark. This probability peaks for a 250 GeV radion and has a branching probability of ~1%. For values as low as $\Lambda_\phi \sim 100$ GeV, this can give a ~20% contribution to the single top production which is about ~80 pb in the SM at the LHC with $\sqrt{s} = 7$ TeV. With the knowledge of the radion mass and with dedicated cuts to isolate the signal from the background this mode can offer a hint to the flavor violating decay of the radion [10]. It will be however be impossible to see any significant effects of flavor violation in the leptonic sector at the LHC in the ATLAS and CMS experiments from radion production. The heavy radion would most likely be seen through its decays to the weak gauge bosons ($m_\phi > 140$ GeV) while the τ mode looks to be significant for the lighter radion. We refer the reader to various detailed studies on radion signals at colliders [6,9].

VIII. CONCLUSIONS

In this paper, we have analyzed the flavor violation in warped extra dimension due to the radion exchange. In this scenario, the SM fermions are propagating in the five-dimensional bulk and the Higgs is localized on the TeV brane. We found that $K - \bar{K}$ and $B_q - \bar{B}_q$ lead to strong constraints on radion mass, m_ϕ and the scale Λ_ϕ. For instance, if $\Lambda_\phi \sim \mathcal{O}(1)$ TeV, one finds that $B_d^0 - \bar{B}_d^0$ implies that $m_\phi \gtrsim 65$ GeV. We have also studied the radion contributions to lepton flavor violating processes: $\ell_i \to \ell_j \phi$ and $\ell_i \to \ell_j \ell_k \ell_l$, in addition to $B \to \ell_i \ell_j$. We have shown that the BR$(\tau \to (e, \mu)\phi)$ imposes a stringent limit on the scale Λ_ϕ for $m_\phi \lesssim \mathcal{O}(1)$ GeV. We emphasized that the radion effect to BR$(B_s \to \mu^+\mu^-)$ can be of order 10^{-8}, which is accessible at the LHCb. We have also analyzed the search for radion at the LHC. Although, we do not find any significant flavor violating signals in the lepton sector, there is definitely a possibility of contributions to single top cross section with the radion decaying through the flavor violating mode of $t\bar{c} + \bar{t}c$.

ACKNOWLEDGMENTS

K. H and A. S acknowledge support from the Academy of Finland under Project No. 137960 and A. S also from the Finnish Cultural Foundation and CIMO. The work of S. K. and A. M. was partially supported by the Science and Technology Development Fund (STDF) under Project No. 1855 and the ICTP under Project No. 30. S. K. R is supported in part by the US Department of Energy under Grant No. DE- FG02-04ER41306.

[1] N. Arkani-Hamed, S. Dimopoulos, and G. R. Dvali, Phys. Lett. B **429**, 263 (1998); I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos, and G. R. Dvali, Phys. Lett. B **436**, 257 (1998); N. Arkani-Hamed, S. Dimopoulos, and G. R. Dvali, Phys. Rev. D **59**, 086004 (1999).
[2] L. Randall and R. Sundrum, Phys. Rev. Lett. **83**, 3370 (1999); **83**, 4690 (1999).
[3] T. Gherghetta and A. Pomarol, Nucl. Phys. **B586**, 141 (2000).
[4] S. J. Huber and Q. Shafi, Phys. Lett. B **498**, 256 (2001).
[5] W. D. Goldberger and M. B. Wise, Phys. Rev. Lett. **83**, 4922 (1999); Phys. Lett. B **475**, 275 (2000).
[6] G. F. Giudice, R. Rattazzi, and J. D. Wells, Nucl. Phys. **B595**, 250 (2001); S. Bae, P. Ko, H. S. Lee, and J. Lee, Phys. Lett. B **487**, 299 (2000); arXiv:hep-ph/0103187; M. Chaichian, A. Datta, K. Huitu, and Z. h. Yu, Phys. Lett. B **524**, 161 (2002).
[7] C. Csaki, M. L. Graesser, and G. D. Kribs, Phys. Rev. D **63**, 065002 (2001); C. Csaki, C. Grojean, J. Hubisz, Y. Shirman, and J. Terning, Phys. Rev. D **70**, 015012 (2004); C. Csaki, J. Hubisz, and S. J. Lee, Phys. Rev. D **76**, 125015 (2007); C. Csaki, C. Delaunay, C. Grojean, and Y. Grossman, J. High Energy Phys. 10 (2008) 055.
[8] K. m. Cheung, Phys. Rev. D **63**, 056007 (2001).
[9] S. R. Choudhury, A. S. Cornell, and G. C. Joshi, arXiv:hep-ph/0012043; U. Mahanta and S. Rakshit, Phys. Lett. B **480**, 176 (2000); U. Mahanta and A. Datta, Phys. Lett. B **483**, 196 (2000); S. C. Park, H. S. Song, and J. H. Song, Phys. Rev. D **65**, 075008 (2002); M. Battaglia, S. De Curtis, A. De Roeck, D. Dominici, and J. F. Gunion, Phys. Lett. B **568**, 92 (2003); A. Datta and K. Huitu, Phys. Lett. B **578**, 376 (2004); P. K. Das, S. K. Rai, and S. Raychaudhuri, Phys. Lett. B **618**, 221 (2005); P. K. Das, Phys. Rev. D **72**, 055009 (2005); M. Toharia, Phys. Rev. D **79**, 015009 (2009); V. P. Goncalves and W. K. Sauter, Phys. Rev. D **82**, 056009 (2010).
[10] A. Azatov, M. Toharia, and L. Zhu, Phys. Rev. D **80**, 031701 (2009).
[11] C. Delaunay, O. Gedalia, S. J. Lee, G. Perez, and E. Ponton, Phys. Rev. D **83**, 115003 (2011).
[12] T. Inami and C. S. Lim, Prog. Theor. Phys. **65**, 297 (1981); **65**, 1772(E) (1981); A. J. Buras, arXiv:hep-ph/0101336.
[13] K. Nakamura *et al.* (Particle Data Group), J. Phys. G **37**, 075021 (2010).
[14] M. Ciuchini *et al.*, J. High Energy Phys. 10 (1998) 008.
[15] G. Buchalla, A. J. Buras, and M. E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996).
[16] A. J. Buras, arXiv:hep-ph/0505175.
[17] L. Lellouch and C. J. D. Lin (UKQCD Collaboration), Phys. Rev. D **64**, 094501 (2001).
[18] D. Becirevic, M. Ciuchini, E. Franco *et al.*, Nucl. Phys. **B634**, 105 (2002).
[19] C. Bobeth, T. Ewerth, F. Kruger, and J. Urban, Phys. Rev. D **64**, 074014 (2001).
[20] H. Davoudiasl and E. Ponton, Phys. Lett. B **680**, 247 (2009).
[21] J. Pumplin, D. R. Stump, J. Huston, H. L. Lai, P. M. Nadolsky, and W. K. Tung, J. High Energy Phys. 07 (2002) 012.