Status of a flavor-maximal nonminimal universal extra dimension model

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In this paper we consider an $S^1/\mathbb{Z}_2$ compactified flat extra dimensional scenario where all the standard model states can access the bulk and have generalized brane localized kinetic terms. The flavor structure of brane kinetic terms for the standard model fermions are dictated by stringent flavor bounds on the first two generations implying an $U(2)_{Q_L} \otimes U(2)_{u_R} \otimes U(2)_{d_R}$ flavor symmetry. We consider the constraints on such a scenario arising from dark matter relic density and direct detection measurements, precision electroweak data, Higgs physics, and LHC dilepton searches. We discuss the possibility of such a scenario providing an explanation of the recently measured anomaly in $R_{K^{(*)}}$ within the allowed region of the parameter space.

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I. INTRODUCTION

Universal extra dimension (UED) [1–3] as an extension of the standard model (SM) of particle physics has received considerable attention in the literature, see [4] and references therein. In these models, one additional spatial dimension with a flat metric is considered compactified on an $S^1/\mathbb{Z}_2$ orbifold. The endpoints where translation symmetry is explicitly broken are locations of two four dimensional space-time hypersurfaces called the 3-branes. The reflection symmetry of the bulk geometry results in a conserved Kaluza-Klein (KK) parity that can stabilize a dark matter (DM) candidate in this setup providing the strongest motivation. In the minimal version of this model (mUED) upper bound on the observed relic density places the scale of new physics at the TeV scale [5,6] implying encouraging prognosis of being explored at collider experiments. Interestingly, the collider phenomenology of these models closely mimics supersymmetric extension of the SM with a relatively compressed spectrum [7].

The radiative corrections modify both the masses and couplings partially lifting the degeneracy in the model [8]. The nonrenormalizability of the 5d theory ensures that the

radiative corrections are proportional to the cutoff and thus incalculable. However a prudent way to accommodate such corrections is to introduce brane localized kinetic terms (BLKT) [9–15]. The BLKT parameters are eventually related to the radiative corrections in a UV complete model, however in this paper we will consider them to be free parameters in the spirit of the so called nonminimal UED models (nmUED). Myriad phenomenological aspects of such a setup has been studied including LHC searches of the strong sector [16,17], Higgs data [18], flavor physics [19–21], unitarity bounds [22], $Z \rightarrow b\bar{b}$ decay width [23], rare top decays [24], and some other sectors [25–29]. If identical BLKT parameters are introduced in the two branes for every bulk field, the KK symmetry is preserved and can lead to a stable DM with a rich phenomenology [30,31].

The BLKT parameters having its origin in the radiative corrections are expected to be dependent on flavor owing to the mass hierarchy in the SM fermions. However, most studies in the literature have remained confined to universal BLKT parameters owing to the dangerous tree level FCNC and level mixing that can arise in the case of most generalized BLKT parameters. Taking a cue from some recent hints about flavor violation beyond the SM in B meson decays [32], in this paper we explore a possible extension of this setup while imposing an $U(2)_{Q_L} \otimes U(2)_{u_R} \otimes U(2)_{d_R}$ flavor symmetry on the BLKT parameters. This ensures the absence of tree level FCNC mediated by SM gauge bosons in the first two generations of the quark sector. Additionally, in the limit where neutrino masses are neglected, the absence of tree level FCNC is ensured in the leptonic sector also [33]. We also keep the BLKT for bosonic degrees of freedom to be identical to prevent level mixing in the gauge sector that can lead to

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nontrivial constraints form the oblique electroweak parameters. Finally we assume a universal brane term for the Yukawas that suppresses flavor violation in the scalar sector while leaving enough freedom in the 5d theory to reproduce the CKM matrix. This setup represents the *flavor maximal nmUED*.

The non-observation of any hint of new physics at the Large Hadron Collider (LHC) either in direct searches for resonances [34,35] or in the increasingly SM like Higgs couplings [36] has been pushing the scale of new physics ever higher. This leads to tension with the upper limit on the scale of the extra dimension coming from the overclosure bound for dark matter relic density. In this paper we explore the possibility of the allowed parameter space in addressing the observed discrepancy in the recent measurement of $R_{K^{(*)}} (= \text{Br}(\bar{B} \rightarrow \bar{K}^{(*)}\mu\mu)/\text{Br}(\bar{B} \rightarrow \bar{K}^{(*)}ee))$ at LHCb [32]. We perform an extensive scan of the parameter space to find the regions of parameter space that is in consonance with the constraints from flavor [32,37–39], dark matter relic density and direct searches [40,41], Higgs data [42], precision electroweak parameters [43], and LHC constraints from the dilepton channel [44,45]. We find only

a very tuned region of parameter space with some large BLKT parameters survive the onslaught.

The rest of the paper is arranged as follows. In the next section we briefly sketch the model and set up the parameter space. In Sec. III we discuss the constraints from the flavor observables $R_{K^{(*)}}$ and $R_{D^{(*)}}$. Dark matter relic density and direct detection observations are discussed in Sec. IV. In Sec. V the phenomenological constraints from Higgs data, oblique parameters, LHC dilepton are briefly reviewed. We then present the results of our extensive parameter scan and discuss their phenomenological implications in Sec. VI before concluding in Sec. VII. We give a few finer details of the model and the relevant flavor violation in the Appendix.

II. THE MAXIMALLY FLAVORED NMUED

We will consider that the extra spatial dimension is compactified on an $S^1/\mathbb{Z}_2$ and all standard model states can access the bulk. The end points are locations of the 3-branes with symmetric brane localized terms. Schematically the five-dimensional action for the SM quarks is given by,

$$\begin{aligned} \mathcal{S}_{\text{quark}} = & \int d^4x \int_0^{\pi R} dy [i\bar{Q}_k \Gamma^M \mathcal{D}_M Q_k + r_{f_k} \{\delta(y) + \delta(y - \pi R)\} i\bar{Q}_k \gamma^\mu \mathcal{D}_\mu P_L Q_k \\ & + i\bar{U}_k \Gamma^M \mathcal{D}_M U_k + r_{f_k} \{\delta(y) + \delta(y - \pi R)\} i\bar{U}_k \gamma^\mu \mathcal{D}_\mu P_R U_k \\ & + i\bar{D}_k \Gamma^M \mathcal{D}_M D_k + r_{f_k} \{\delta(y) + \delta(y - \pi R)\} i\bar{D}_k \gamma^\mu \mathcal{D}_\mu P_R D_k], \end{aligned} \quad (1)$$

where the subscript k represents the flavor index. We follow similar prescription for leptons. Further details of the model including the KK decomposition leading to the KK towers for all bulk fields, is given in Appendix A.

The masses of the KK excitation for SM quarks, represented by M_{Q_n} , can be obtained by solving the following transcendental equations,

$$r_f M_{Q_n} = \begin{cases} -2 \tan\left(\frac{M_{Q_n} \pi R}{2}\right) & \forall n \text{ even,} \\ 2 \cot\left(\frac{M_{Q_n} \pi R}{2}\right) & \forall n \text{ odd.} \end{cases} \quad (2)$$

Evidently, for $r_f = 0$ this reduces to the UED KK-mass n/R .

Following the notation of [24], the gauge, scalar, and Yukawa sector Lagrangians are given by,

$$\begin{aligned} \mathcal{S}_{\text{gauge}} = & -\frac{1}{4} \int d^4x \int_0^{\pi R} dy \\ & \times \left[\sum_a (\mathcal{F}^{MNa} \mathcal{F}_{MN}^a + r_g \{\delta(y) + \delta(y - \pi R)\} \mathcal{F}^{\mu\nu a} \mathcal{F}_{\mu\nu}^a) \right. \\ & \left. + \mathcal{B}^{MN} \mathcal{B}_{MN} + r_g \{\delta(y) + \delta(y - \pi R)\} \mathcal{B}^{\mu\nu} \mathcal{B}_{\mu\nu} \right], \end{aligned} \quad (3)$$

$$\begin{aligned} \mathcal{S}_{\text{scalar}} = & \int d^4x \int_0^{\pi R} dy [(\mathcal{D}^M \Phi)^\dagger (\mathcal{D}_M \Phi) + \tilde{\mu}_h^2 \Phi^\dagger \Phi - \tilde{\lambda}_h (\Phi^\dagger \Phi)^2 \\ & + \{\delta(y) + \delta(y - \pi R)\} (r_\phi (\mathcal{D}^\mu \Phi)^\dagger (\mathcal{D}_\mu \Phi) \\ & + \mu_B^2 \Phi^\dagger \Phi - \lambda_B (\Phi^\dagger \Phi)^2)], \end{aligned} \quad (4)$$

$$\begin{aligned} \mathcal{S}_{\text{Yuk}} = & - \int d^4x \int_0^{\pi R} dy [\tilde{y}_{ij}^u \bar{Q}_i \tilde{\Phi} U_j + \tilde{y}_{ij}^d \bar{Q}_i \Phi D_j \\ & + r_y \{\delta(y) + \delta(y - \pi R)\} \\ & \times (\tilde{y}_{ij}^u \bar{Q}_{iL} \tilde{\Phi} U_{jR} + \tilde{y}_{ij}^d \bar{Q}_{iL} \Phi D_{jR}) + \text{H.c.}]. \end{aligned} \quad (5)$$

In this study we will assume universal BLKT parameter for electroweak gauge and scalar sectors, i.e., $r_g = r_\phi$ which simplifies the gauge sector [17,46,47]. Moreover, in the fermionic sector we will consider flavor dependent BLKT parameters modulo a global flavor symmetry of $U(2)_{Q_L} \otimes U(2)_{u_R} \otimes U(2)_{d_R}$. This suppresses dangerous flavor changing neutral (FCNC) and charged currents (FCCC) in the quark sector which is highly constrained from various flavor observations, e.g., $K - \bar{K}$ oscillation, ϵ -parameter etc. [48]. The Yukawa couplings explicitly

TABLE I. Choice of BLT parameters.

$(\nu_e)_L, e_R$	r_{f_1}	$(u)_L, u_R, d_R$	r_{f_1}
$(\nu_\mu)_L, \mu_R$	r_{f_2}	$(c)_L, c_R, s_R$	r_{f_1}
$(\nu_\tau)_L, \tau_R$	r_{f_3}	$(t)_L, t_R, b_R$	r_{f_3}
W_μ, B_μ, G_μ, H	r_g	Yukawa	r_y

break this symmetry leading to an acceptable minimal flavor violation framework [49]. We consider $r_g > r_{f_i}$ to ensure that the lightest KK particle (LKP) is the $B_\mu^{(1)}$. The BLKT setup used in the phenomenological studies in this paper has been summarized in the Table I. We will present all our results in terms of the scaled dimensionless BLKT variables R_X where $R_X = r_X/R$. For $R_X < -\pi$ the masses obtained from Eq. (2) becomes imaginary giving rise to tachyonic zero modes [16,17]. Throughout the paper we will restrict ourselves to ranges for the BLKT parameters from 0 to 8. The choice of the BLKT parameters, are explicitly presented in Table I.

In the KK parity conserving scenario the even KK mode gauge bosons have couplings with the SM fermions, i.e., the zero mode fermions. In the gauge basis the effective Lagrangian can be written as,

$$\mathcal{L}_X = \sum_{X=\gamma,Z} \frac{X_\mu^{(n)}}{2c_w} [g_{f_L}^{X(n)} \bar{f}_L \gamma^\mu f_{iL} + g_{f_R}^{X(n)} \bar{f}_R \gamma^\mu f_{iR}], \quad (6)$$

where c_w is the usual cosine of the Weinberg angle, and the couplings are explicitly given as,

$$g_{f_L}^{Z(n)} = 2(T_{3f} - Q_f g s_w^2) I^{(n)}, \quad g_{f_R}^{Z(n)} = -2Q_f g s_w^2 I^{(n)}, \quad (7a)$$

$$g_{f_{L,R}}^{\gamma(n)} = 2Q_f g s_w c_w I^{(n)}, \quad (7b)$$

Note that since we are not taking different BLKT parameters for left- and right-handed fermions $\gamma^{(n)}$ will have the same coupling with both. The overlap integral $I^{(n)}$ is given by,

$$I^{(n)} = \frac{\sqrt{R_g + \pi}}{R_{f_i} + \pi} \sqrt{\frac{2}{\pi}} \frac{R_{f_i} - R_g}{\sqrt{1 + \frac{R_g^2 M_{gn}^2}{4} + \frac{R_g}{\pi}}} \quad (8)$$

Evidently, in Eqs. (6) and (7), the KK mode n is even due to KK parity conservation. The interaction between KK modes of W -boson and SM fermions are given by,

$$\mathcal{L}_W = \frac{1}{\sqrt{2}} W_\mu^{(n)} [g_{d_L}^W \bar{u}_j P_L \gamma^\mu d_j + g_{\ell_L}^W \bar{\nu}_j P_L \gamma^\mu \nu_j], \quad (9)$$

where $g_{d_L}^W$ and $g_{\ell_L}^W$ are given by $gI(R_g, R_{f_j}, n)$, which are actually the overlap integral times the gauge couplings. Note that we take the BLKT parameters for the same generation of quarks and leptons to be the same.

The Lagrangians in Eqs. (6) and (9) are in the gauge basis. In the mass basis this leads to FCNC and FCCC, mediated by the KK gauge bosons. This is because of the matrix

$$G^{(n)} = \text{diag}(g_{f_1}^{X(n)/W^{(n)}}, g_{f_2}^{X(n)/W^{(n)}}, g_{f_3}^{X(n)/W^{(n)}}),$$

being diagonal but not proportional to the identity matrix. However, note that the SM GIM mechanism would still be in effect because $G^{(0)}$ is still proportional to the identity matrix owing to the orthonormality of the mode functions. As far as the leptonic sector is concerned, KK gauge boson mediated FCNC is proportional to the neutrino Yukawa couplings. In this paper we will neglect the neutrino masses and thus avoiding the FCNC in the leptonic sector. Owing to the different BLKT parameters for the lepton generations the strength of the KK gauge boson couplings to the SM charged leptons would be dependent on the generational index. In the Appendix B we give the details of the flavor violation in the gauge and Yukawa sector.

III. FLAVOR OBSERVABLES $R_{K^{(*)}}$ AND $R_{D^{(*)}}$

Recent measurement of the flavor observables $R_{K^{(*)}}$ and $R_{D^{(*)}}$ shows a consistent $\sim 2\sigma$ discrepancy with the SM. This has stimulated a plethora of models both from bottom-up [50–56] and top-down approach [57–63] to explain this tension. In this section we briefly discuss the impact of these observables as well as $R_{D^{(*)}}$ measurements on the parameter space of the model.

The current LHCb measurements of $R_{K^{(*)}}$ are given by [32,64],

$$R_{K^{(*)}} = \frac{\text{Br}(\bar{B} \rightarrow \bar{K}^{(*)} \mu \mu)}{\text{Br}(\bar{B} \rightarrow \bar{K}^{(*)} e e)}, \quad (10a)$$

$$R_K = 0.745_{-0.074}^{+0.090} \pm 0.036, \quad R_{K^*} = 0.69_{-0.07}^{+0.11} \pm 0.05. \quad (10b)$$

Integrating out the heavy KK modes from the interaction terms presented in Eq. (6) we can obtain the relevant four fermion operators which contribute in the flavor observable $R_{K^{(*)}}$ which is basically a $\Delta F = 1$ transition. Following the standard definitions of Wilson coefficients $C_{9,10}^{(\ell)}$ [65] we can write the new physics contributions ($\Delta C_{9,10}^{(\ell)}$) in our model as,

$$\Delta C_9 = -\sum_{\substack{X=\gamma,Z \\ \{n\}}} \frac{c g_{\mu\nu}^{X(n)}}{M_{X(n)}^2} (g_{b_L}^{X(n)} - g_{s_L}^{X(n)}),$$

$$\Delta C'_9 = -\sum_{\substack{X=\gamma,Z \\ \{n\}}} \frac{c g_{\mu\nu}^{X(n)}}{M_{X(n)}^2} (g_{b_R}^{X(n)} - g_{s_R}^{X(n)}), \quad (11a)$$

$$\Delta C_{10} = \sum_{\substack{X=\gamma,Z \\ \{n\}}} \frac{c g_{\mu A}^{X(n)}}{M_{X(n)}^2} (g_{b_L}^{X(n)} - g_{s_L}^{X(n)}),$$

$$\Delta C'_{10} = \sum_{\substack{X=\gamma,Z \\ \{n\}}} \frac{c g_{\mu A}^{X(n)}}{M_{X(n)}^2} (g_{b_R}^{X(n)} - g_{s_R}^{X(n)}), \quad (11b)$$

where the common factor $c = \pi/(2\sqrt{2}G_F\alpha c_w^2)$, and $g_{f_{V,A}}^{X(n)} = (g_{f_L}^{X(n)} \pm g_{f_R}^{X(n)})/2$. In these equations the sum over even KK modes n is also taken. Now, in terms of these Wilson coefficients R_K is given by [66],

$$R_K = \frac{\sum_{j=9,10} |C_j^{\text{SM}} + \Delta C_j^\mu + \Delta C_j^{\prime\mu}|^2}{\sum_{j=9,10} |C_j^{\text{SM}} + \Delta C_j^e + \Delta C_j^{\prime e}|^2}, \quad (12)$$

where $C_9^{\text{SM}} \simeq -C_{10}^{\text{SM}} \simeq 4.2$ [33]. Also, following [60,67] R_{K^*} can be presented as,

$$\frac{R_{K^*}}{R_K} = 1 + \frac{1}{(C^{\text{SM}})^2 R_K} [2C^{\text{SM}}\{\mathcal{A}^\mu - (\mu \rightarrow e)\} + \{|\mathcal{A}^\mu|^2 + |\mathcal{B}^\mu|^2 - (\mu \rightarrow e)\}], \quad (13)$$

where $C^{\text{SM}} = C_9^{\text{SM}} + C_{10}^{\text{SM}}$, and $\mathcal{A}^\mu = \Delta C_9^\mu + \Delta C_{10}^{\prime\mu} - \Delta C_{10}^\mu$; $\mathcal{B}^\mu = \Delta C_9^\mu + \Delta C_{10}^\mu + \Delta C_{10}^{\prime\mu}$. Within the experimental uncertainties, this ratio can be evaluated to be $R_{K^*}/R_K = 0.94 \pm 0.18$ [60].

We generate BLKT parameters randomly in the range (0,8) and $1/R$ in the range (500,2500) GeV and calculate R_K using Eq. (12) and R_{K^*}/R_K using Eq. (13) for each set. The resulting values have been plotted in Fig. 2. Within the left-right symmetric BLKT framework (i.e., the BLKT parameters for left- and right-handed fermions are identical) considered here the primed Wilson coefficients [see Eq. (11)] are not vanishing. However, we obtain large number of parameter points that simultaneously satisfy 1σ (2σ) bounds from R_K and R_{K^*} .

We now consider the flavor observables $R_{D^{(*)}}$ in this model. The standard definition of which is

$$R_{D^{(*)}} = \frac{\text{Br}(\bar{B} \rightarrow D^{(*)} \tau \nu_\tau)}{\text{Br}(\bar{B} \rightarrow D^{(*)} \ell \nu_\ell)}, \quad (\ell = e \text{ or } \mu). \quad (14)$$

The world average of $R_{D^{(*)}}$ from experimental results from BABAR, Belle, and LHCb [37–39] is given by [68],

FIG. 1. Schematic Feynman diagram for the process (a) $b \rightarrow s\mu\bar{\mu}(e\bar{e})$ and (b) $b \rightarrow cl\bar{\nu}_l$.

$$R_D = 0.407 \pm 0.039 \pm 0.024,$$

$$R_{D^*} = 0.304 \pm 0.013 \pm 0.007. \quad (15)$$

The contribution to $R_{D^{(*)}}$ comes after integrating out the KK modes of W -boson. The effective operator can be written as,

$$\mathcal{L}_W^{\text{eff}} = -\frac{4G_F}{\sqrt{2}} V_{cb} \sum_{\ell} C_\ell^{(n)} (\bar{c} \gamma_\mu P_L b) (\bar{\ell} \gamma^\mu \nu_\ell), \quad (16)$$

where the KK index n is even and the Wilson coefficients $C_\ell^{(n)}$ are given by [33],

$$C_\ell^{(n)} = \left(\frac{m_W}{m_{W^{(n)}}} \right)^2 g'_{\mu_L} g'_{\ell_L}. \quad (17)$$

In terms of the Wilson coefficients defined in Eqs. (16) and (17), $R_{D^{(*)}}$ can be written as [33,69],

$$R_{D^{(*)}}(C_\tau, C_\mu) = 2R_{D^{(*)}}^{\text{SM}} \frac{|1 + C_\tau|^2}{1 + |1 + C_\mu|^2}, \quad (18)$$

where the SM values are given by [70–72],

$$R_D^{\text{SM}} = 0.300 \pm 0.008, \quad R_{D^*}^{\text{SM}} = 0.252 \pm 0.003.$$

Unlike $b \rightarrow s$ transition, the $b \rightarrow c$ transition can occur in the SM, albeit CKM suppressed, in tree level, i.e., $n = 0$ is allowed in the right panel of Fig. 1. Thus the nmUED contribution of higher KK mode W -bosons in the $b \rightarrow c$ process, and hence in $R_{D^{(*)}}$, is quite small which is evident from the Fig. 3 where not much deviation is found in $R_{D^{(*)}}$ even after considering BLKT values up to 20. This corroborates similar observations made in the context of warped extra dimensional setup [62]. In the rest of the paper, we ignore the observables $R_{D^{(*)}}$.

IV. DARK MATTER

The preservation of KK parity due to the equal-strength BLKT parameters at the two boundary points implies the

FIG. 2. Allowed parameter points in the (a) $R_K - 1/R$ and (b) $R_{K^*}/R_K - 1/R$ plane. The black hatched region corresponds to 1σ deviation from experimental results and the region enclosed by the uppermost and lowermost black lines corresponds to 2σ deviation from experimental results. In (b), blue points denote those allowed by 2σ of R_K and orange points denote those allowed by 1σ of R_K .

FIG. 3. Allowed parameter points from flavor observables (a) R_D and (b) R_{D^*} . The black hatched region corresponds to 1σ deviation from experimental results and the region enclosed by the uppermost and lowermost black lines corresponds to 2σ deviation from experimental results.

stability of the lightest KK particle (LKP). Since in the nmUED setup the BLKT parameters determine the mass spectrum as well as the couplings, the identity of the LKP can vary depending on the choice of BLKT parameters. In the mUED it's always the first KK level photon $A^{(1)}$, which is generally denoted in the literature by $B^{(1)}$ since this is the dominant part in $A^{(1)}$ [5,6,73,74]. In nmUED, however, $Z^{(1)}$ (i.e., $W_3^{(1)}$), $H^{(1)}$ can also be viable LKPs apart from the usual $A^{(1)}$. The detailed studies of dark matter in the nmUED has been performed in [30,31]. In this paper we will stick to the case where $A^{(1)}$ is the DM candidate which is mainly the $B^{(1)}$. Since for a specific value of the compactification radius, the KK masses decrease with increasing values of BLKT parameters, we have to take $r_g > r_{f_i}$ to avoid fermion LKP. Also in our analysis we do not take into account the co-annihilation processes involving $W_3^{(1)}$, $H^{(1)}$ etc. Actually we stick to the conservative overclosure bound only when considering the parameter

space, i.e., we demand the relic density of the $A^{(1)}$ to be just less than the observed relic density of the universe.

The relic density of $A^{(1)}$ can be written, by solving the corresponding Boltzmann equation with appropriate assumptions, as

$$\Omega_{A^{(1)}} h^2 \approx \frac{1.04 \times 10^9}{M_{\text{Pl}}} \frac{x_F}{\sqrt{g_*(x_F)}} \frac{1}{a + \frac{3b}{x_F}}, \quad (19)$$

where M_{Pl} is the Planck mass, $x_F (= m_{A^{(1)}}/T_F)$ represents the freeze-out temperature which we take ~ 25 [5,6], the quantities a and b are the coefficients of the nonrelativistic expansion of the thermally averaged annihilation cross section, i.e., $\langle \sigma v \rangle = a + b \langle v^2 \rangle + \mathcal{O}(\langle v^4 \rangle) \approx a + 6b/x$ and lastly the quantity g_* is the total number of effective degrees of freedom.

The main annihilation channels of $A^{(1)}$ are $A^{(1)}A^{(1)} \rightarrow f\bar{f}$ and $A^{(1)}A^{(1)} \rightarrow HH^*$, where f represents the SM fermions

FIG. 4. (a) Allowed parameter points in the relic density, $\Omega_{A^{(1)}} h^2$ vs $1/R$ plane. The black line corresponds to the overclosure bound 0.12 [75]. (b) The LKP $A^{(1)}$ dependence on the relic density. The overclosure bound fixes the mass of the LKP $M_{A^{(1)}} \sim 600$ GeV.

and H is the SM Higgs. The thermally averaged cross section of these annihilation channels can be expanded in terms of the coefficients a and b which via Eq. (19) gives the estimation of relic density. In our analysis we also considered the coannihilation channels with gauge or scalar particles in the initial states, such channels are: (i) $A^{(1)} W_3^{(1)} \rightarrow f \bar{f}$, (ii) $A^{(1)} W_3^{(1)} \rightarrow H H^*$, (iii) $A^{(1)} W^{(1)\pm} \rightarrow f \bar{f}'$, (iv) $A^{(1)} G^{(1)} \rightarrow W^\mp G^\pm$, (v) $A^{(1)} G^{(1)\pm} \rightarrow Z(\gamma) G^\pm$. We have used the expressions for cross sections of these processes from [6].

The DM direct detection searches can put stringent constraints on the parameter space of the model. This can even be more crucial given the ever-increasing precision of modern direct detection experiments like, LUX, Xenon1T etc. [40,41]. The scattering cross section of the DM $A^{(1)}$ off nuclei is ultimately related to its scattering from quarks. In the nonrelativistic limit the total cross section can have both the spin-independent and spin-dependent parts. The details of the analysis of the direct detection constraints are discussed in [30,31].

We generate BLKT parameters randomly in the range (0,8) and $1/R$ in the range (500,2000) GeV and calculate $\Omega_{A^{(1)}} h^2$ using Eq. (19) for each set. In Fig. 4 we show the allowed points which are also allowed by direct detection constraints from LUX (spin-dependent) [40] and Xenon1T (spin-independent) [41]. We observe that a large number of parameter points remain allowed by the overclosure bound. However, it sets strict upper limits of $1/R$ at ~ 1.6 TeV and $M_{A^{(1)}}$ at ~ 600 GeV.

V. PHENOMENOLOGICAL CONSTRAINTS

A. Higgs data

The production of Higgs boson at the LHC dominantly occurs through the gluon fusion process $gg \rightarrow H$ which is driven by fermion triangle loops, and within the SM the top quark loop dominates due to its large Yukawa coupling with the Higgs. The subsequent important decay of the SM

Higgs to the diphoton is also a loop induced process where SM fermions and W -boson run in the triangle loop. In the nmUED, the KK modes of SM fermions as well as W -boson can contribute in these loop induced processes and can potentially modify the effective couplings. One can thus put stringent constraints on the model parameter space from the observed Higgs couplings [18,76–80]. The cross section and decay width for the Higgs including the KK contributions can be written as,

$$\sigma_{gg \rightarrow H} = \frac{G_F \alpha_s^2 m_H^3}{16\sqrt{2}} \left| \frac{1}{2} A_t(\tau_t) + C(r_{\text{NP}}) \mathcal{A}_{\text{NP}}(\tau_{\text{NP}}) \right|^2, \quad (20)$$

$$\Gamma_{H \rightarrow \gamma\gamma} = \frac{G_F \alpha^2 m_H^3}{128\sqrt{2}} \left| A_W(\tau_W) + 3 \left(\frac{2}{3} \right)^2 A_t(\tau_t) + N_{c,\text{NP}} Q_{\text{NP}}^2 \mathcal{A}_{\text{NP}}(\tau_{\text{NP}}) \right|^2, \quad (21)$$

where $\tau_j = m_H^2/4m_j^2$, $C(r_{\text{NP}})$ is the $SU(3)$ color factor, $N_{c,\text{NP}}$ is the number of color states of the new physics (NP) particles, Q_{NP} is the electric charge of the NP particle in the loop. In the light Higgs limit ($m_H \ll 2m_j$) $A_t \sim 4/3$, $A_W \sim -7$. The quantity $\mathcal{A}_{\text{NP}}$ is defined as,

$$\mathcal{A}_{\text{NP}}(\tau_{\text{NP}}) = v \frac{\partial}{\partial v} \log [\det(\mathcal{M}(v))] A(\tau_{\text{NP}}), \quad (22)$$

where $\mathcal{M}(v)$ is the mass matrix for respective particles. Clearly, $\mathcal{A}_{\text{NP}}$ contains only contribution from top quark for the $\sigma_{gg \rightarrow H}$ but both top quark and W boson for $\Gamma_{H \rightarrow \gamma\gamma}$.

With these one can now calculate the dimensionless parameters $c_{gg} = \sigma_{gg \rightarrow h}^{\text{NP}} / \sigma_{gg \rightarrow h}^{\text{SM}}$ and $c_{\gamma\gamma} = \Gamma_{gg \rightarrow h}^{\text{NP}} / \Gamma_{gg \rightarrow h}^{\text{SM}}$ to compare the modification induced by the KK modes with the experimentally observed values. Using the values obtained in [42], we have $\sqrt{c_{gg}} = 0.88 \pm 0.11$ and $\sqrt{c_{\gamma\gamma}} = 1.18 \pm 0.12$. In our analysis we will use this constraint.

TABLE II. Details of the choice of parameter range for numerical scans. Set-II represents a refined scan around BP-III defined in Table IV.

	No. of points (in million)	R_{f_1}	R_{f_2}	R_{f_3}	R_g	$1/R$ (GeV)
Set-I	2.0	(0, 8)	(0, 8)	(0, 8)	(0, 8)	(500, 2500)
Set-II	0.28	(5.5, 6.6)	(4.2, 5.5)	(0.0, 1.0)	(7.3, 8.5)	(1300, 1500)

TABLE III. Survival fraction after subsequent imposition of various constraints.

Total	$R_K^{2\sigma}, (R_{K^*}/R_K)^{2\sigma}$	DM	Higgs	EWPO	Dilepton
1	0.2	0.01	2×10^{-4}	4×10^{-5}	2×10^{-6}

Total	$R_K^{1\sigma}, (R_{K^*}/R_K)^{1\sigma}$	DM	Higgs	EWPO	Dilepton
1	0.03	1×10^{-3}	4×10^{-5}	2×10^{-6}	0

B. Oblique parameters

The oblique corrections to the electroweak gauge boson propagators, incarnated in the Peskin-Takeuchi parameters a.k.a., S , T , and U parameters, generally put strong constraints on the BSM models. In nmUED these electroweak constraints are discussed in [9,17,19,24,26]. Here we briefly discuss them for completeness.

In KK parity conserving nmUED the Fermi constant G_F receives corrections due to the presence of tree-level coupling between second KK level gauge boson and the SM fermions. This correction can be presented as,

$$G_F = G_F^0 + \delta G_F = \frac{g^2}{4\sqrt{2}M_W^2} + \sum_{\substack{k \geq 2 \\ k \in \text{even}}} \frac{(gI^{(k)})^2}{4\sqrt{2}M_{W^{(k)}}^2}, \quad (23)$$

where G_F^0 (δG_F) represents the s -channel SM (even KK mode) $W^\pm$ -boson exchange, and $I^{(k)}$ is the relevant overlap integral shown in Eq. (8). With these quantities the nonzero Peskin-Takeuchi parameters can be written as [17,19,24],

$$T_{\text{nmUED}} = -\frac{1}{\alpha} \frac{\delta G_F}{G_F}, \quad U_{\text{nmUED}} = \frac{4 \sin^2 \theta_W}{\alpha} \frac{\delta G_F}{G_F}. \quad (24)$$

The most recent fit to the electroweak precision data given by Gfitter group [43],

$$S = 0.05 \pm 0.11, \quad T = 0.09 \pm 0.13, \quad U = 0.01 \pm 0.11. \quad (25)$$

We can now perform the model parameter scan following the prescription of [24] by defining the χ^2 in terms of the covariance matrix. For a maximal 2σ (3σ) deviation, we need $\chi^2 \leq 6.18$ (9.21) given the two degrees of freedom.

C. Dilepton bounds

The presence of KK number violating (but KK parity conserving) interactions can result in the single production of second KK excitation of gauge bosons. Resonant production of the second level KK gauge boson and its subsequent decay to SM fermions at the LHC, i.e., $pp \rightarrow X^{(2)} \rightarrow l^+ l^-$ proceeds via couplings of the form $g_{X^{(2)} f \bar{f}}$.

FIG. 5. Allowed parameter points in the $M_{A^{(1)}} - M_{q_3^{(1)}}$ plane. The blue points represent the points allowed by (a) 1σ and (b) 2σ values of R_K and R_{K^*}/R_K . The green points show the allowed points after additionally imposing the DM constraints i.e., the overclosure bound and direct detection cross sections. The red points are allowed after additionally imposing the Higgs decay width constraints. Finally the black points represent the allowed points after additionally imposing the electroweak precision constraints.

TABLE IV. The three benchmark points that are allowed by all the constraints.

	R_g	R_{f_1}	R_{f_2}	R_{f_3}	$M_{A^{(1)}} \text{ (GeV)}$	$M_{q_3^{(1)}} \text{ (GeV)}$	R_K	R_{K^*}/R_K	Ωh^2
BP-I	5.098	3.922	2.966	0.2385	540.8	1108	0.924	1.034	0.101
BP-II	6.404	4.997	4.087	0.0281	539.1	1295	0.9358	1.028	0.101
BP-III	7.788	6.102	4.746	0.3496	525.8	1249	0.9279	1.032	0.0958

The dilepton searches at the LHC can put stringent constraints on these second level KK gauge bosons. Recent ATLAS [44] and CMS [45] searches have been used to put constraint on the parameter space of nmUED in [31]. Without taking recourse to the extensive collider simulations we simply translate the bounds from [31] on our simulated sample points. We compare the coupling $g_{X^{(2)}f\bar{f}}$ in our model with the constraint given in [31] for a given $M_{X^{(2)}}$ to determine the exclusion from LHC dilepton searches. Recasting other collider studies [16,17] of nmUED is beyond the scope of this work. However, we expect enough freedom exists within the framework to evade these constraints due to the presence of nonuniversal BLKT parameters.

VI. CONSTRAINTS ON THE PARAMETER SPACE

In this section we discuss the constraints on the parameter space obtained by imposing various phenomenological constraints elaborated in the preceding sections.

We perform a simulation of two million randomly generated parameter points with BLKT parameters in the range (0,8) and $1/R$ in the range (500, 2500) GeV. In Table II, the Set-I represents the full set of randomly generated parameter points. For each set of input parameters all the observables are calculated. The various constraints are sequentially imposed on this generated sample points. The fractional cutflow for surviving points after implementation of subsequent constraints is depicted in Table III. We present the points allowed by R_K , R_{K^*}/R_K , dark matter, Higgs and EWPT constraints in the $M_{A^{(1)}}-M_{q_3^{(1)}}$ plane in Fig. 5. We are left with just three points from our sample set that are within 2σ of R_K and R_{K^*}/R_K , and survive the additional imposition of dilepton constraints. These are shown as benchmark points in Table IV. No point survives all constraints that is simultaneously within the 1σ band of R_K and R_{K^*}/R_K . This has to be contrasted with Fig. 2(b) where we obtain a large number of points which are within the 1σ band of the two flavor observables. This clearly espouses the tension between the flavor observables and precision observables, Higgs data, and dark matter observations. We conclude from this generic scan that the surviving points are tuned and accord a marginal improvement in $R_{K^{(*)}}$.

We identify BP-III of Table IV as a point relatively insulated from collider constraint and perform a finer scan

around this point defined by Set-II of Table II. We show the allowed points in the parameter space of $M_{A^{(1)}} - M_{q_3^{(1)}}$ in Fig. 6. These points represent a section of the parameter space that remains allowed by all constraints and can explain the $R_{K^{(*)}}$ observation within 2σ .

A qualitative analysis of the cuts imposed points towards the region of the parameter space that has survived. The arguments are as follows:

- (1) The tree level flavor violation in the quark sector [Fig. 1(a)] is proportional to $|R_{f_2} - R_{f_3}|$. Thus a sizeable flavor violation indicated by R_{K^*} prefers a considerable difference between the values of R_{f_2} and R_{f_3} .
- (2) In Eqs. (11) the vectorial Wilson coefficient $|\Delta C_9|$ dominates over $|\Delta C_{10}|$. From Eq. (12) it is evident that for $R_K < 1$, one needs $|\Delta C_9^\mu|$ to be greater than $|\Delta C_9^e|$. And since $|\Delta C_9^e|$ is inversely proportional to R_{f_i} , we require R_{f_1} to be greater than R_{f_2} .
- (3) The increasingly SM like Higgs data tunes R_{f_3} to small positive values.
- (4) The allowed parameter space from precision data closes down the difference between R_{f_2} and R_ϕ as evident from Eqs. (23) and (8).

As we have assumed R_ϕ to be greater than R_{f_i} to get a bosonic LKP, following the above arguments we conclude the following hierarchy structure of the BLKT

FIG. 6. Allowed parameter points which are within 2σ of $R_{K^{(*)}}$ and are allowed by other constraints are presented in the parameter space of $M_{A^{(1)}} - M_{q_3^{(1)}}$.

parameters: $R_\phi \geq R_{f_1} \geq R_{f_2} > R_{f_3}$, which is exactly the distribution we have found in our benchmark points.

VII. SUMMARY AND CONCLUSION

Nonminimal version of the UED has a rich phenomenology owing to the presence of boundary localized parameters which play a consequential role in determining not only the mass spectrum of the KK particles but also significant deviations in the relevant couplings. In this paper we present a variant of nmUED with nonuniversal BLKT parameters with an $U(2)_{Q_L} \otimes U(2)_{u_R} \otimes U(2)_{d_R}$ symmetry in the fermionic sector.

This flavor maximal setup implies new flavor violating contributions mediated by the KK gauge bosons while contributions from the Yukawa sector remain suppressed by the Yukawa couplings. Utilizing this flavor violation we explore the possibility of reconciling the recent discrepancy between the observation and SM expectation in flavor observables $R_{K^{(*)}}$ and $R_{D^{(*)}}$ while conforming to other collider and cosmological constraints. We perform an extensive scan of the parameter space to find the allowed regions that is in consonance with all constraints. We present an updated study of the dark matter aspects of the model by using the relic density observations from Planck and latest direct detection cross section results from LUX and Xenon1T which already constrains the $B^{(1)}$ LKP mass in the ballpark of 600 GeV. The other observables which we consider include (i) Higgs data, namely the $H \rightarrow \gamma\gamma$ and $gg \rightarrow H$ channels where the effect of KK modes will be most dominant; (ii) electroweak precision observables a.k.a. oblique parameters; (iii) bounds from dilepton searches of ATLAS and CMS.

We find that while the new contributions are significant in bridging the gap for the $R_{K^{(*)}}$, the improvements in terms of $R_{D^{(*)}}$ is numerically insignificant. The extensive scan reveals that even with all the flexibility afforded by nonuniversality of the BLKT parameters only a fine-tuned region of parameter space is capable of evading the phenomenological constraints while remaining within 2σ of the measured values of R_K and R_{K^*}/R_K . While we find no parameter point that is within 1σ .

As nonobservation of new physics at the LHC pushes the KK scale ever higher it is in imminent danger of being severely constrained from the upper bound on the LKP arising from the dark matter relic density overclosure bound. It is evident from the analysis presented in this paper that the requirements to match the flavor discrepancies are in tension with constraints from precision physics, Higgs data, and dark matter observations. Relaxation of this tension requires further generalization of the BLKT framework. For example, one might consider different BLKT parameter for chiral components of the same flavor or explicit violation of KK parity by introducing different BLKT parameter at different branes. A detailed and

rigorous study of the collider, flavor, and cosmological aspects of such a generalized framework would be in order.

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APPENDIX A: MODEL DETAILS

The five-dimensional four component quark fields (Q , U , D) can be represented as their corresponding Kaluza-Klein (KK) excitations as,

$$Q(x, y) = N_Q^{(0)} Q_L^{(0)} + \sum_{n=1}^{\infty} [Q_L^{(n)}(x) f_L^{(n)}(y) + Q_R^{(n)}(x) g_L^{(n)}(y)], \quad (\text{A1a})$$

$$U(x, y) = N_Q^{(0)} U_R^{(0)} + \sum_{n=1}^{\infty} [U_L^{(n)}(x) f_R^{(n)}(y) + U_R^{(n)}(x) g_R^{(n)}(y)], \quad (\text{A1b})$$

$$D(x, y) = N_Q^{(0)} D_R^{(0)} + \sum_{n=1}^{\infty} [D_L^{(n)}(x) f_R^{(n)}(y) + D_R^{(n)}(x) g_R^{(n)}(y)]. \quad (\text{A1c})$$

In the effective 4D theory the zero modes of Q generates the $SU(2)_L$ doublet quarks and the zero modes of U (D) are identified with the up (down) type singlet quarks, i.e., after compactification and orbifolding the zero modes of Q gives the left-handed doublet comprising of SM t_L and b_L , whereas t_R and b_R emerges from the U and D respectively. The compact form of quark doublet is $Q \equiv (Q_i, Q_j)^T$, where i and j correspond to up type and down type quark respectively. $N_Q^{(0)}$ is the normalization constant of the fermionic wave functions for zero-mode. In Eq. (1) the terms containing the parameter r_f are the boundary localized kinetic terms (BLKTs). In the mUED, r_f is assumed to be vanishing, thus by setting the BLKT parameters to zero one can translate from nmUED to mUED. Also note that the Latin indices in Eq. (1) run from 0 to 4 whereas the Greek indices from 0 to 3. We will be using the metric convention $g_{MN} \equiv \text{diag}(+1, -1, -1, -1, -1)$. The covariant derivative, $D_M \equiv \partial_M - i\tilde{g}W_M^a T^a - i\tilde{g}'B_M Y$, where $\tilde{g}$ and $\tilde{g}'$ are the 5D gauge

coupling constants of $SU(2)_L$ and $U(1)_Y$, respectively, and T^a and Y are the corresponding generators. The five-dimensional gamma matrices are $\Gamma^M = (\gamma^\mu, -i\gamma_5)$.

The variation of the action with appropriate boundary conditions lead to the y -dependent mode functions f and g as,

$$f_L(y) = g_R(y) = N_Q^{(n)} \begin{cases} \frac{\cos[M_{Q_n}(y - \frac{\pi R}{2})]}{C_{Q_n}} & \forall n \text{ even,} \\ -\frac{\sin[M_{Q_n}(y - \frac{\pi R}{2})]}{S_{Q_n}} & \forall n \text{ odd,} \end{cases} \quad (\text{A2})$$

and

$$g_L(y) = f_R(y) = N_Q^{(n)} \begin{cases} \frac{\sin[M_{Q_n}(y - \frac{\pi R}{2})]}{C_{Q_n}} & \forall n \text{ even,} \\ \frac{\cos[M_{Q_n}(y - \frac{\pi R}{2})]}{S_{Q_n}} & \forall n \text{ odd,} \end{cases} \quad (\text{A3})$$

with

$$C_{Q_n} = \cos\left(\frac{M_{Q_n} \pi R}{2}\right), \quad S_{Q_n} = \sin\left(\frac{M_{Q_n} \pi R}{2}\right). \quad (\text{A4})$$

The mode functions f and g satisfy the orthonormality conditions,

$$\begin{aligned} \int dy [1 + r_f \{\delta(y) + \delta(y - \pi R)\}] k^{(m)}(y) k^{(n)}(y) \\ = \delta^{mn} = \int dy l^{(m)}(y) l^{(n)}(y) \end{aligned} \quad (\text{A5})$$

where, k can be f_L or g_R and l corresponds to g_L or f_R . From the above condition one can obtain the normalization factors as

$$N_{Q_n} = \sqrt{\frac{2}{\pi R}} \left[\frac{1}{\sqrt{1 + \frac{r_f^2 M_{Q_n}^2}{4} + \frac{r_f}{\pi R}}} \right]. \quad (\text{A6})$$

For zero-mode, the normalization constant is given by

$$N_{Q0} = \frac{1}{\sqrt{r_f + \pi R}}. \quad (\text{A7})$$

Note that $r_f = 0$ implies the usual (m)UED normalization $\sqrt{2/(\pi R)}$ for n th mode, whereas for zero-mode it is $\sqrt{1/(\pi R)}$. The y -profile solutions for the gauge field V^μ (with $V = W, Z, \gamma$) is given by the Eqs. (A2) with M_{Q_n} being replaced by M_{Φ_n} . The M_{Φ_n} is the solution of transcendental equations [Eq. (2)] for gauge fields with BLKT parameter r_ϕ instead of r_f . The orthonormality condition will be similar to that of $k^{(i)}(y)$ of Eq. (A5). The corresponding normalization factor is given by $N_{\Phi 0} = 1/\sqrt{r_\phi + \pi R}$.

APPENDIX B: FLAVOR VIOLATION: GAUGE AND YUKAWA SECTOR

Here we give a brief outline of the flavor violation that can arise in the model for different choices of model parameters.

In the gauge sector the relevant interactions of $W^{(n)\pm}$ and $Z^{(n)}$ are shown in Eqs. (6) and (9). After EWSB, the quark mass matrices are diagonalized by unitary matrices V_u and V_d which are the constituents of the CKM matrix $V_{\text{CKM}} = V_u^\dagger V_d$. Now, flavor violation occurs both in the neutral and charged currents in the mass eigenstate by $(V_{d_{L,R}}^\dagger \mathcal{I}^{(n)} V_{d_{L,R}})$ and $V_{u_L}^\dagger \mathcal{I}^{(n)} V_{d_L}$ respectively, where $\mathcal{I}^{(n)} = \text{diag}\{I^{(n)}(R_g, R_{f_1}), I^{(n)}(R_g, R_{f_2}), I^{(n)}(R_g, R_{f_3})\}$. Evidently, the flavor violation is generated due to the fact that the matrix $\mathcal{I}^{(n)}$ is diagonal but not proportional to identity matrix. This is mandated by the generation-wise different BLT parameter, i.e., $r_{f_1} \neq r_{f_2} \neq r_{f_3}$.

The Yukawa Lagrangian for the charged scalar part in the presence of the boundary localized Yukawa term can be written as,

$$\begin{aligned} \mathcal{L}_{\text{Yuk}}^{\text{charged}} = \int d^4x \int_0^{\pi R} dy [1 + r_y \{\delta(y) + \delta(y - \pi R)\}] \\ \times (-\tilde{y}_d \bar{Q}_L \phi^{(n)+} D). \end{aligned} \quad (\text{B1})$$

Here we write the interactions for the down sector; the discussion for the up sector follows similarly. Also, we do not take any flavor structure in r_y itself. The 5D Yukawa $\tilde{y}_d$ and its 4D counterpart are related by the relation,

$$\tilde{y}_{d_i} = \frac{y_{d_i}}{\pi R} (r_{f_i} + \pi R) \sqrt{r_\phi + \pi R}, \quad (\text{B2})$$

which can be obtained by taking $r_y = 0$ and considering the zero mode wave functions of corresponding fields with appropriate normalization factors. Now, to check the flavor violation in the Yukawa sector, we consider the coupling between the higher KK mode charged scalar and the zero mode fermions. This coupling can be obtained from Eq. (B1). By taking appropriate normalizations into account one can show this coupling to be,

$$-y_{d_i} \frac{\sqrt{r_\phi + \pi R}}{\pi R} \int_0^{\pi R} dy f_\phi^{(2)}(y). \quad (\text{B3})$$

Since this coupling is r_{f_i} independent there will be no flavor violation.

We can also take $r_f \neq 0 \neq r_y$, where r_y is same for all three generations. In that case, the 5D Yukawa $\tilde{y}_d$ coupling is related to its 4D counterpart by,

$$\tilde{y}_{d_i} = y_{d_i} \frac{(r_{f_i} + \pi R)}{(r_y + \pi R)} \sqrt{r_\phi + \pi R}. \quad (\text{B4})$$

For the above condition, the Eq. (B3) can be rewritten as

$$-y_{d_i} \frac{\sqrt{r_\phi + \pi R}}{r_y + \pi R} \int_0^{\pi R} dy [1 + r_y \{\delta(y) + \delta(y - \pi R)\}] f_\Phi^{(2)}(y) \quad (\text{B5})$$

which is again independent of r_f , so there is no flavor violation. But if we take r_y different for all generations or set $r_{y_i} = r_{f_i}$, then we will get flavor violations in the Yukawa sector due to the presence of the terms like $1/(r_{y_i} + \pi R)$ or $1/(r_{f_i} + \pi R)$ respectively in the $\mathcal{I}^{(n)}$

matrix. In this article, we actually consider the second one, i.e., $r_f \neq 0 \neq r_y$ and r_y being same for all generations, we do not have any flavor violation in the scalar sector.

In passing we mention that for the BLKT parameter choice, $r_y \neq r_{f_i}$ there exists KK level mixing in the Yukawa interactions. Due to the KK parity conservation this mixing occurs between even-even modes or odd-odd modes. The mixing angle between the n th and $(n + 2k)$ th ($k \in \mathbb{Z}$) KK level can be estimated as $\theta_{\text{mix}} \sim I(r_y, r_{f_i}, n, k) m_{f_i} R / 2k$, where $I(r_y, r_{f_i}, n, k)$ is the corresponding overlap integral. This mixing angle is suppressed by the new physics scale, see [16,18] for details.

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