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B Meson Anomalies in a Pati-Salam Model within the Randall-Sundrum Background

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Lepton number as a fourth color is the intriguing theoretical idea of the famous Pati-Salam (PS) model. While in conventional PS models, the symmetry breaking scale and the mass of the resulting vector leptoquark are stringently constrained by $K_L \to \mu e$ and $K \to \pi\mu e$, the scale can be lowered to a few TeV by implementing the PS gauge group in the five-dimensional Randall-Sundrum background. The symmetry breaking is achieved by the choice of boundary conditions and the constraints from kaon physics can be avoided by implementing the SM fermions as bulk fields with different zero mode localizations for the three generations. We consider the flavour phenomenology of this model in the context of the hints for lepton flavour universality violation in semileptonic B decays. Concerning $b \to s\ell^+\ell^-$ transitions, the observed deviations from the SM predictions (including $R(K)$ and $R(K^*)$) can be explained with natural values for the free parameters of the model without fine-tuning in other flavour observables (like $B_s - \overline{B}_s$ mixing). Even though we find sizable effects in $R(D)$, $R(D^*)$ and $R(J/\Psi)$ one cannot account for the current central values in the constrained setup of our minimal model due to the stringent constraints from $D - \bar{D}$ mixing and $\tau \to 3\mu$.

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I. INTRODUCTION

So far, the Large Hadron Collider (LHC) at CERN did not directly observe any particles beyond the ones of the Standard Model (SM) of particle physics. However, we have accumulated intriguing hints for lepton flavor universality (LFU) violation in semi-leptonic B decays within recent years. Most prominently, there exist deviations from the SM predictions in $b \to s\mu^+\mu^-$ above the $5\,\sigma$ level [1] [69] and the combination of the ratios $R(D)$ and $R(D^*)$ differs by $4.1\,\sigma$ from its SM prediction [2]. Furthermore, also $R(J/\Psi)$ points towards the violation of LFU in $b \to c\tau\nu$ processes [3]. This suggests a possible connection between these two classes of decays and motivates the investigation of simultaneous explanations [4–23].

In fact, the $SU(2)$ singlet vector leptoquark (VLQ) with hypercharge $2/3$ is a natural candidate for a simultaneous explanation [5, 6, 9]. It contributes to $b \to s\mu^+\mu^-$ as well as to $b \to c\tau\nu$ and it does not couple down-quarks to neutrinos, avoiding the bounds from $B \to K^{(*)}\nu\nu$. This allows for large flavour violating effects and the bounds from direct searches [24] and EW precision data [25, 26] can be avoided [17, 20]. Interestingly, this LQ appears in the theoretically very appealing Pati-Salam (PS) [27] model and several attempts have been made in the literature to construct a model addressing the flavour anomalies based on the corresponding gauge symmetry [21–23, 28].

In conventional PS models, the bounds on the breaking scale from $K_L \to \mu e$ and $K \to \pi\mu e$ are so strong (at the PeV scale) [29, 30] that any other observable effect in flavour physics is ruled out. Nonetheless, if the PS gauge symmetry is implemented in the 5D Randall-Sundrum (RS) background [31], the mass scale of the Kaluza-Klein (KK) resonances (including the VLQ) can be much lower, i.e. in the few TeV range [70]. The suppression of the lepton flavour violating kaon decays can be achieved by introducing the SM fermions as zero modes of bulk fermions [32] with their couplings to the KK modes determined by their localization along the RS bulk. Since the zero mode localizations are free parameters, one can obtain the required non-trivial flavour structure in order to give interesting effects in $b \to s\mu^+\mu^-$ and $b \to c\tau\nu$ transitions.

II. THE MODEL

Our starting point is a 5D RS space-time [31]

$$ds^2 = e^{-2ky}\eta_{\mu\nu}dx^\mu dx^\nu - dy^2\,, \qquad 0 \le y \le \pi R \qquad (1)$$

with the PS [27] bulk gauge symmetry $SU(4)\times SU(2)_L \times SU(2)_R$. The symmetry is broken to its SM subgroup by means of boundary conditions on the UV brane. Note that the unbroken $U(1)_Y$ is a linear combination of $U(1)_{B-L}$ (contained in $SU(4)$) and $U(1)_R$ of $SU(2)_R$. Therefore, relaxing the assumption of a discrete left-right

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symmetry, one can always account for the measured values of g_Y

As in the conventional PS model, the SM fermions are embedded into complete representations of the PS gauge group. In addition, they are introduced as bulk fields in the RS background and the zero modes correspond to the SM quarks and leptons. Their localizations are determined by their 5D bulk masses [32], but can be altered by the presence of brane-kinetic terms. Since on the UV brane only the SM gauge symmetry is unbroken, the localizations of quarks and leptons of the same generation can differ from each other [71]. The Higgs doublet is introduced as a 4D field confined to the UV brane, hence its couplings to the KK modes are strongly suppressed. This choice ensures the compliance with electroweak precision constraints.

The 4D dual theory, according to the AdS/CFT correspondence [33], is a composite model with a global PS symmetry and KK modes corresponding to composite resonances. The gauging of the SM subgroup explicitly breaks the global symmetry in the elementary sector. Hence, the SM fermions are partially composite due to a linear mixing of the elementary fermions with composite operators of the same quantum numbers. Therefore, the simplified version of our model, according to the deconstruction approach [34], contains composite vector resonances of the $SU(4)\times SU(2)_L\times SU(2)_R$ symmetry group with common mass M, as well as three generations of heavy vector-like quarks and leptons corresponding to the first KK modes.

Concerning fermions we have three generations of chiral (SM) fermions, the quark doublets q_i, the lepton doublets ℓ_i, the quark singlets d_i and u_i as well as the lepton singlet e_i (and the right-handed neutrino which we do not consider in the following as it is not relevant for our discussion). In addition, we have the three generations of vector-like fermions which we denote by the corresponding capital letters. The mass terms for the fermion before electroweak symmetry breaking read

$$\begin{aligned}\mathcal{L}_M \;=\; & -M^L_{ij}\left(\bar{Q}^L_i Q^R_j + \bar{L}^L_i L^R_j\right)\\ & -M^R_{ij}\left(\bar{U}^L_i U^R_j + \bar{D}^L_i D^R_j + \bar{E}^L_i E^R_j\right)\\ & -m^{qL}_{ij}\bar{q}_i Q^R_j - m^{\ell L}_{ij}\bar{\ell}_i L^R_j\\ & -m^{uR}_{ij}\bar{U}^L_i u_j - m^{dR}_{ij}\bar{D}^L_i d_j - m^{\ell R}_{ij}\bar{E}^L_i e_j + h.c.\,.\end{aligned} \tag{2}$$

Here the superscripts L and R denote the chirality of the vector-like fields. Note that Q_i, L_i, U_i, D_i and E_i, are embedded into complete representations under the PS gauge group, enforcing equality of the respective mass terms. Without loss of generality, we can work in a basis where M^L_{ij} and M^R_{ij} are diagonal in flavour space, and to a good approximation, the masses of the composite states are universal:

$$M^L_{ij}\approx M^R_{ij}\approx M\delta_{ij}\,. \tag{3}$$

For the terms mixing vector-like fermions with the SM ones we assume for simplicity the absence of mixing with the right-handed SM $SU(2)$ singlets, i.e. $m^{fR}_{ij}=0$. In RS models without brane-kinetic terms m^{fL}_{ij} and M^L_{ij} are diagonal in the same basis. Assuming that the brane-kinetic terms are also diagonal in that basis, one can write

$$m^{fL}=\begin{pmatrix} M^f_1 & 0 & 0\\ 0 & M^f_2 & 0\\ 0 & 0 & M^f_3\end{pmatrix}. \tag{4}$$

In the following, we will assume M_1 to be zero or negligibly small [72]. Therefore, the first generation is purely elementary while the second and third generation of left-handed SM quarks and leptons are partially composite. Since all terms are flavour diagonal the problem reduces to diagonalizing several 2×2 matrices mixing SM with vector-like fermions. For the quarks and leptons we achieve this by the transformation

$$\begin{pmatrix} f^i_L\\ F^1_L\end{pmatrix}\to\begin{pmatrix} c^f_i & -s^f_i\\ s^f_i & c^f_i\end{pmatrix}\begin{pmatrix} f^i_L\\ F^i_L\end{pmatrix} \tag{5}$$

with $i=2,3$, $f=q,\ell$, $F=Q,L$, and $s^f_i=\sin\alpha^f_i$ with $\alpha^f_i=\arctan\left(M^f_i/M\right)$.

After EW symmetry breaking additional mass terms for the fermions, originating from the Yukawa couplings, arise. Now, the 3×3 sub-block of the light fermions is in general not diagonal in the same basis as M^{fL}_{ij} and m^{fL}_{ij}. In order to avoid tree-level flavour changing neutral currents (FCNCs) in the down quark sector, we assume the down Yukawa coupling to be aligned with m^{qL} and M^{qL}, i.e. diagonal in the same basis. In the left-handed up-quark sector, tree-level FCNCs are then unavoidable, but are determined and suppressed by small off-diagonal CKM elements.

With these assumptions, the couplings of KK gauge bosons to fermions are given by

$$L^V_{ff}=i\bar{f}_i\gamma_\mu\left(g^{V*}_L\Gamma^{VL}_{f_if_j}P_L+g^{V*}_R\Gamma^{VR}_{f_if_j}P_R\right)f_jV^\mu\,. \tag{6}$$

Here, $V=g$, $W^\pm$, W^3, B$-$L, LQ with the coupling $g^{V*}_{L,R}=\theta g^V_{\rm SM}$ being enhanced by the RS volume $\theta=\sqrt{k\pi R}\sim 6$ with respect to the elementary gauge coupling of the gauge boson $V_{\rm SM}$. For the LQ and the B-L gauge boson $g^{\rm LQ*}_L=\theta\frac{g_s}{\sqrt{2}}$ and $g^{\rm B-L*}_L=\theta\frac{\sqrt{3}}{2\sqrt{2}}g_s$, respectively. The couplings of the KK modes of $SU(2)_R$ to the SM fermions are small.

The relevant matrices in flavour space $\Gamma^{L,V}_{ij}$ read:

$$\Gamma^{V^0,L}_{d_id_j}=\begin{pmatrix}0&0&0\\0&s^{q2}_2&0\\0&0&s^{q2}_3\end{pmatrix}_{ij} \tag{7}$$

$$\Gamma^{V^0,L}_{u_iu_j}=V^{\rm CKM}_{ik}\begin{pmatrix}0&0&0\\0&s^{q2}_2&0\\0&0&s^{q2}_3\end{pmatrix}_{kl}V^{\rm CKM\,*}_{jl} \tag{8}$$

$$\Gamma^{LQ,L}_{d_i\ell_j} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & s_2^q s_2^\ell c_\ell & s_2^q s_2^\ell s_\ell \\ 0 & -s_3^q s_3^\ell s_\ell & s_3^q s_3^\ell c_\ell \end{pmatrix}_{ij}, \tag{9}$$

$$\Gamma^{W,L}_{u_i d_j} = V^{\rm CKM}_{ik} \begin{pmatrix} 0 & 0 & 0 \\ 0 & s_2^{q2} & 0 \\ 0 & 0 & s_3^{q2} \end{pmatrix}_{kj}, \tag{10}$$

$$\Gamma^{V,L}_{\ell_i\ell_j} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & s_2^{\ell 2} c_\ell^2 + s_3^{\ell 2} s_\ell^2 & \left(s_2^{\ell 2} - s_3^{\ell 2}\right) s_\ell c_\ell \\ 0 & \left(s_2^{\ell 2} - s_3^{\ell 2}\right) s_\ell c_\ell & s_3^{\ell 2} c_\ell^2 + s_2^{\ell 2} s_\ell^2 \end{pmatrix}_{ij} \tag{11}$$

Here we neglected flavour mixing with the first generation and dropped the flavour-universal θ^2-suppressed terms [34]. In this limit $\Gamma^{R,V}_{ij} = 0$. $V^{\rm CKM}$ denotes the CKM matrix, V^0 stands for the electrically neutral gauge bosons and s_ℓ parametrizes the misalignment in flaovur space between $m^{\ell L}$, $M^{\ell L}$ and the lepton Yukawa coupling in the 2-3 sector.

III. OBSERVABLES

A. $R(D)$ and $R(D^*)$

We define the effective Hamiltonian for $b \to c\ell\nu$ transitions as

$$H^{\ell_f\nu_i}_{\rm eff} = \frac{4G_F}{\sqrt{2}} V_{cb} C_L^{fi} \left[\bar{c}\gamma^\mu P_L b\right] \left[\bar{\ell}_f \gamma_\mu P_L \nu_i\right], \tag{12}$$

where in the SM $C_L^{fi} = \delta_{fi}$ and the contribution of our model is given by

$$C_L^{fi} = \frac{\sqrt{2}}{4G_F V_{cb}} \frac{\kappa^*_{3f} V_{2k} \kappa_{ki}}{M^2} + \theta \Gamma^{W,L}_{\ell_f\ell_i} \frac{m_W^2}{M^2}. \tag{13}$$

Here the first term originates from the LQ, with $\kappa_{ij} = \theta \frac{g_s}{\sqrt{2}} \Gamma^{{\rm LQ},L}_{d_i\ell_j}$, while the second term is due to the KK mode of the $W^\pm$. Thus we find

$$R(X)/R(X)_{\rm SM} = |1 + C_L^{33}|^2 + \sum_{i=1}^{2} |C_L^{3i}|^2, \tag{14}$$

with $X = D, D^*, J/\Psi$.

This has to be compared to the experimental measurements of $R(D)$, $R(D^*)$ and $R(J/\Psi)$ [3]. A global fit assuming NP in C_L only gives [35]

$$C_L^{\rm NP} = 0.131 \pm 0.033. \tag{15}$$

B. $b \to s\ell^+\ell^-$ transitions

Using the effective Hamiltonian

$$H^{\ell_f\ell_i}_{\rm eff} = -\frac{4G_F}{\sqrt{2}} V_{tb} V^*_{ts} \sum_{a=9,10} C_a^{fi} O_a^{fi},$$

$$O^{fi}_{9(10)} = \frac{\alpha}{4\pi} \left[\bar{s}\gamma^\mu P_L b\right] \left[\bar{\ell}_f \gamma_\mu (\gamma^5) \ell_i\right], \tag{16}$$

we have

$$C_9^{fi} = -C_{10}^{fi} = \frac{-\sqrt{2}}{2G_F V_{tb} V^*_{ts}} \frac{\pi}{\alpha} \frac{\kappa_{2i}\kappa^*_{3f}}{M^2}. \tag{17}$$

The allowed $2\,\sigma$ range is given by [1]

$$-0.37 \geq C_9^{22} = -C_{10}^{22} \geq -0.88. \tag{18}$$

Concerning lepton flavour violating B decays, we use the results of Ref. [36] for the analysis of $B \to K^{(*)}\tau\mu$. The only experimental limit for $\mu\tau$ final states is [37]

$${\rm Br}\left[B \to K\tau\mu\right]_{\rm EXP} \leq 4.8 \times 10^{-5}, \tag{19}$$

at 90% confidence level, and the corresponding prediction for our case of $C_9 = -C_{10}$ reads

$${\rm Br}\left[B \to K\tau\mu\right] = 1.96 \times 10^{-8} \left(\left|C_9^{23}\right|^2 + \left|C_9^{32}\right|^2\right). \tag{20}$$

C. $D - \bar{D}$ mixing

$D^0 - \bar{D}^0$ mixing receives tree level contributions from the KK modes of the gluon, the $B - L$ gauge boson and the W^3. The resulting NP contribution to the effective Hamiltonian $H_{\rm eff} = C_L Q_1 + h.c.$ is

$$C_L = \frac{\theta^2 \left(\frac{3}{4}g_s^2 + \frac{1}{2}g_2^2\right)}{4M^2} \left(V_{cs} V^*_{us} s_2^{q2} + V_{cb} V^*_{ub} s_3^{q2}\right)^2, \tag{21}$$

with $Q_1 = (\bar{c}\gamma^\mu P_L u)(\bar{c}\gamma_\mu P_L u)$. We have for the matrix element

$$M_{12}^D = \frac{1}{3} m_D f_D^2 B_1^D(\mu) \eta_D(\mu) C_L, \tag{22}$$

with $B_1^D(3\,{\rm GeV}) \approx 0.76$ [38], $\eta_D(3{\rm GeV}) = 0.77$ [39, 40] and $f_D \approx (0.212)\,{\rm GeV}$ [41, 42]. Using the HFLAV results of CKM 2016 [2], the imaginary part of the matrix element should satisfy

$$|{\rm Im}[M_{12}]| < 2 \times 10^{-16}\,{\rm GeV}. \tag{23}$$

D. $\tau \to 3\mu$

The neutral B-L and W^3 KK gauge bosons mediate the decay $\tau \to 3\mu$. Using the results of [36] and neglecting contributions suppressed by g_Y/θ, we find

$${\rm Br}\left[\tau \to 3\mu\right] = \frac{m_\tau^5 \tau_\tau}{768\pi^3} \frac{1}{M^4} \left| \sum_V^{W^3,{\rm B\text{--}L}} g_L^{V*2} \Gamma^{VL}_{\ell_2\ell_3} \Gamma^{VL}_{\ell_2\ell_2} \right|^2. \tag{24}$$

Here τ_τ is the tau lifetime. This result has to be compared to the current experimental bound of 1.2×10^{-8} [43].

E. $B_s - \overline{B}_s$ mixing

Due to the assumed flavour alignment in the left-handed down-quark sector, our model does not only forbid tree level contributions to $B_s - \overline{B}_s$ mixing, but also makes the one-loop contributions to $B_s - \overline{B}_s$ mixing finite (even in unitary gauge) due to a GIM-like mechanism. In addition, flavour violation solely originates from Yukawa couplings. Thus, the effect is very efficiently suppressed by $1/M^4$.

F. Direct LHC searches

The most stringent constraints on the KK mass scale stem from direct LHC searches for resonances decaying to $t\bar{t}$, dijet or $\tau\bar{\tau}$. $t\bar{t}$ resonance searches constrain the RS KK gluon mass to be above 3.3 TeV in the case of bulk fermions and flavor anarchy [44]. In our setup, however, the branching ratio into $t\bar{t}$ final states is significantly smaller than in the flavour-anarchic scenario, so that we can conservatively lower the mass scale of the lightest resonances to $M = 3.0\,\mathrm{TeV}$. Due to the reduced branching ratio into $t\bar{t}$, the dijet final state is relevant in our setup. The most recent CMS constraint on heavy dijet resonances [45] is nonetheless still weaker than the aforementioned $t\bar{t}$ constraint. Both the B-L gauge boson and the W^3 KK mode contribute to the $\tau\bar{\tau}$ final state. Comparing the Z' of the sequential SM, for which [46] finds $M_{Z'} > 2.42\,\mathrm{TeV}$, with our model, we find that the larger branching ratio into the $\tau\bar{\tau}$ final state is counteracted by a significantly reduced production cross section: first generation quarks do not couple to the B-L gauge boson, and their coupling to the W^3 KK mode is suppressed by $1/\theta$. The $t\bar{t}$ resonance constraint of $\approx 3\,\mathrm{TeV}$ is hence the strongest limit on the KK mass scale M.

IV. PHENOMENOLOGY

Since we aim at getting a large effect in $b \to c\tau\nu$ transitions, a large compositeness of the third generation is required. In addition, M should not be too large and we therefore use a mass of 3 TeV. In order to get a sizeable effect in $b \to s\mu^+\mu^-$, while not violating the upper limit on the $\tau \to 3\mu$ branching ratio, moderate values of s_2^ℓ are preferred. In the left plot of Fig. 1 we therefore show the allowed regions in the s_2^q-s_ℓ plane for $s_2^\ell = 0.2$, $s_3^\ell = 1/\sqrt{2}$ and $s_3^q = \sqrt{3}/2$. At this benchmark point $R(X)/R(X)_{\mathrm{SM}} \approx 1.07$. One can see that $b \to s\mu^+\mu^-$ can be explained at the $1\,\sigma$ level without violating bounds from $D - \bar{D}$ mixing or $\tau \to 3\mu$.

In the right plot of Fig. 1 we show the correlations between $R(X)/R(X)_{\mathrm{SM}}$ and $b \to s\mu^+\mu^-$ by scanning over s_3^q, s_2^q, s_3^ℓ, s_2^ℓ and s_ℓ. Only the parameter points consistent with all experimental bounds are shown. We see that in general a large effect in $b \to s\mu^+\mu^-$ limits the size of the possible effect in $R(X)/R(X)_{\mathrm{SM}}$ and vice versa. Furthermore, the solution of the $b \to s\mu^+\mu^-$ anomaly in our model predicts a large branching ratio for $\tau \to 3\mu$ within the reach of Belle II.

Due to the constraints from $D^0 - \bar{D}^0$ mixing and $\tau \to 3\mu$, we do not obtain sizeable effects neither in $b \to s\tau^+\mu^-$ nor $\tau \to \phi\mu$ [15], nor in $b \to s\tau^+\tau^-$ transitions as recently examined in Ref. [47].

V. CONCLUSIONS AND OUTLOOK

In this article we considered a PS model embedded in the RS space-time in which the symmetry is broken down to the SM one by boundary conditions on the endpoints of the extra dimension. While in previous models based on the PS symmetry the effect in $b \to c\tau\nu$ was only generated by the vector LQ, we have as well a W' contribution which enhances in our setup the total NP effect in $b \to c\tau\nu$ processes by roughly 80%. Still, we find that one cannot fully account for $b \to c\tau\nu$ data due to the stringent constraints from $D - \bar{D}$ mixing. However, an $O(5\%)$ effect in $R(X)/R(X)_{\mathrm{SM}}$ is possible. Furthermore, the model can naturally explain the anomaly in $b \to s\mu^+\mu^-$ transitions including the hints for the violation of lepton flavour universality from $R(K)$ and $R(K^*)$. In addition, our model predicts small effect in $b \to s\tau\mu$ and $b \to s\tau\tau$ transitions, while the effect in $\tau \to 3\mu$ is sizable and also the CP violation in the $D - \bar{D}$ system is close to the current experimental values.

In our minimal setup we assumed right-handed fermions and the Higgs to be elementary. Giving up these assumptions, one obtains an even richer phenomenology and also an explanation of the tensions in the anomalous magnetic moment of the muon [48, 49] and/or in ε'/ε [50] could become possible.

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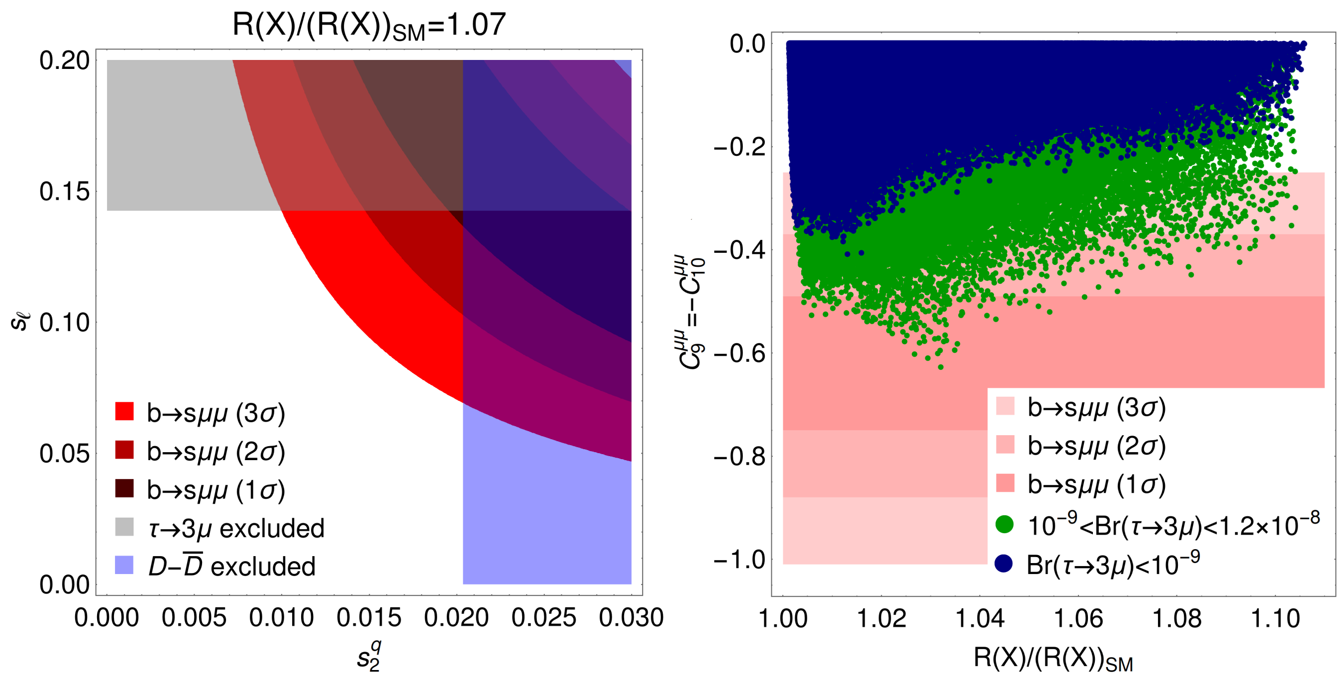

FIG. 1: Left: Allowed regions from $b \to s\mu^+\mu^-$ (red) and the exclusion limits from $D-\bar{D}$ mixing (blue) and $\tau \to 3\mu$ (gray) for $M = 3\,\text{TeV}$, $s_2^\ell = 0.2$, $s_3^\ell = 1/\sqrt{2}$ and $s_3^q = \sqrt{3}/2$. With these values, $R(X)/R(X)_{\rm SM} \approx 1.07$ (with $X = D, D^*, J/\Psi$). We see that $b \to s\mu^+\mu^-$ can be explained at the $1\,\sigma$ level without violating the bounds from other observables. Right: Correlations between $R(X)/R(X)_{\rm SM}$ and $b \to s\mu^+\mu^-$ for $M = 3\,\text{TeV}$. Here we scanned over $0.3 < s_3^q < \sqrt{3}/2$, $0 < s_2^q < 0.2$, $0.3 < s_3^\ell < \sqrt{3}/2$, $0 < s_2^\ell < 0.2$ and $0 < s_\ell < 0.3$. Only the parameter points consistent with $D-\bar{D}$ mixing are shown. As we can see, the predicted branching ratio for $\tau \to 3\mu$ is large and very well within the reach of Belle II.

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[69] Including only $R(K)$ and $R(K^*)$ the significance is at the $4\,\sigma$ level [51–56].
[70] The B physics anomalies have been considered in the context of ordinary composite Higgs or extra dimension models in Refs. [57–63]. For general flavor analyses of conventional RS models, see Refs. [64–68].
[71] We thank Csaba Csáki for reminding us of this possibility.
[72] This is motivated by the fact that in order to explain the anomalies no couplings to electrons are required and that once couplings to electrons are present, there arise stringent bounds from $\mu \to e\gamma$ and $\mu \to 3e$.