Minimal Unified Resolution to $R_{K^{(*)}}$ and $R(D^{(*)})$ Anomalies with Lepton Mixing

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(Received 30 June 2017; published 12 October 2017)

It is a challenging task to explain, in terms of a simple and compelling new physics scenario, the intriguing discrepancies between the standard model expectations and the data for the neutral-current observables R_K and R_{K^*}, as well as the charged-current observables $R(D)$ and $R(D^*)$. We show that this can be achieved in an effective theory with only two unknown parameters. In addition, this class of models predicts some interesting signatures in the context of both B decays as well as high-energy collisions.

Introduction and the data.—Several recent hints of discrepancies in a few charged- as well as neutral-current semileptonic decays of B mesons have intrigued the community. Unlike the case for fully hadronic decay modes that suffer from large (and, in some cases, not-so-well-understood) strong interaction corrections, the theoretical uncertainties in semileptonic decays are much better controlled. Even these uncertainties are removed to a great extent in ratios of similar observables. While, individually, none of the observables militate against the standard model (SM), viewed together, they strongly suggest that some new physics (NP) is lurking around the corner [1,2]. The pattern also argues convincingly for the violation of lepton-flavor universality.

With the ratios of partial widths being particularly clean probes of physics beyond the SM, on account of the cancellation of the leading uncertainties, let us focus on $R(D)$ and $R(D^*)$ defined as

$$R(D^{(*)}) \equiv \frac{\mathrm{BR}(B \to D^{(*)}\tau\nu)}{\mathrm{BR}(B \to D^{(*)}\ell\nu)}, \qquad \ell \in \{e,\mu\} \quad (1)$$

and analogous ratios for the neutral-current sector

$$R_{K^{(*)}} \equiv \frac{\mathrm{BR}(B \to K^{(*)}\mu\mu)}{\mathrm{BR}(B \to K^{(*)}ee)}. \quad (2)$$

With the major source of uncertainty in the individual modes being the form factors, they largely cancel out in ratios like $R(D^{(*)})$ or $R_{K^{(*)}}$, and the SM estimates for these ratios are rather robust. [The cancellation works best for relatively large momentum transfers (where the leptonic mass effects are negligible), the region with the best data.] Several measurements of $R(D)$ and $R(D^*)$ by the *BABAR* [3], Belle [4,5], and LHCb [6,7] Collaborations indicated an upward deviation from the SM expectations. Combining the individual results, namely, $R(D) = 0.407 \pm 0.039 \pm 0.024$ and $R(D^*) = 0.304 \pm 0.019 \pm 0.029$, the

discrepancies are at $\sim 2.3\sigma$ and $\sim 3.4\sigma$, respectively. On the inclusion of the correlation between the data, the combined significance is at the $\sim 4.1\sigma$ level [8] from the SM predictions [9].

The data on R_K and R_{K^*}, on the other hand, lie systematically below the SM expectations [10,11]:

$$R_K = 0.745^{+0.090}_{-0.074} \pm 0.036, \qquad q^2 \in [1\!:\!6]\ \mathrm{GeV}^2,$$
$$R_{K^*}^{\mathrm{low}} = 0.660^{+0.110}_{-0.070} \pm 0.024, \qquad q^2 \in [0.045\!:\!1.1]\ \mathrm{GeV}^2,$$
$$R_{K^*}^{\mathrm{cntr}} = 0.685^{+0.113}_{-0.069} \pm 0.047, \qquad q^2 \in [1.1\!:\!6]\ \mathrm{GeV}^2. \quad (3)$$

For both R_K and $R_{K^*}^{\mathrm{cntr}}$, the SM predictions are virtually indistinguishable from unity [12], whereas for $R_{K^*}^{\mathrm{low}}$ it is ~ 0.9 (owing to a finite m_μ). Except for $R_{K^*}^{\mathrm{low}}$, the theoretical uncertainties have been subsumed in the experimental ones. Thus, the measurements of R_K, $R_{K^*}^{\mathrm{low}}$, and $R_{K^*}^{\mathrm{cntr}}$, respectively, correspond to 2.6σ, 2.1σ, and 2.4σ shortfalls from the SM expectations.

For the K^* mode, a discrepancy is visible not only in the ratios of binned differential distribution for muon and electron modes but also in some angular distributions, like the celebrated P_5' [13] anomaly for the decay $B \to K^*\mu\mu$ [14], at more than 3σ. Restricting ourselves to only the low- and medium-q^2 region, namely, $q^2 \leq 6\ \mathrm{GeV}^2$ (as the high-q^2 region can be affected by a different kind of physics [15]), we do not include this anomaly in our analysis. However, we see later that our fitted Wilson coefficients can explain this discrepancy as pointed out in global fits [1].

A similar suppression (at a level of approximately 3σ) is seen in the observable $\Phi \equiv d\mathrm{BR}(B_s \to \phi\mu\mu)/dm_{\mu\mu}^2$ in the analogous bin ($m_{\mu\mu}^2 \in [1\!:\!6]\ \mathrm{GeV}^2$) [16–18], namely,

$$\Phi = \begin{cases} (2.58^{+0.33}_{-0.31} \pm 0.08 \pm 0.19) \times 10^{-8}\ \mathrm{GeV}^{-2} & (\mathrm{exp}), \\ (4.81 \pm 0.56) \times 10^{-8}\ \mathrm{GeV}^{-2} & (\mathrm{SM}). \end{cases}$$
$$(4)$$

With a low theoretical error, this bin is virtually the same as that for R_K and $R_{K^*}^{\rm cntr}$. This suggests strongly that the discrepancies in the latter have been caused by a depletion of the $b \to s\mu\mu$ channel rather than an enhancement in $b \to see$, a surmise further vindicated by the P'_5 anomaly. Note that P'_5 is dominated by the vector operator $\mathcal{O}_9$, while the two-body decay $B_s \to \mu\mu$ is controlled by the axial vector operator $\mathcal{O}_{10}$, both of them defined later.

With possible corrections from large $\Delta\Gamma_s$, as well as next-to-leading-order (NLO) electroweak and next-to-next-to-leading-order QCD corrections calculated, the SM prediction is quite robust with only small uncertainties accruing from the Cabibbo-Kobayashi-Maskawa matrix elements and the decay constant of B_s. The LHCb measurement at a significance of 7.8σ [19,20] shows an excellent agreement between the data and the SM:

$$\mathrm{BR}(B_s \to \mu\mu) = \begin{cases} (3.0 \pm 0.6^{+0.3}_{-0.2}) \times 10^{-9} & (\mathrm{exp}), \\ (3.65 \pm 0.23) \times 10^{-9} & (\mathrm{SM}), \end{cases} \quad (5)$$

and hence puts very strong constraints on NP models, in particular, on those incorporating (pseudo)scalar or axial-vector currents [21]. However, note that the central value can accommodate an $\sim 20\%$ suppression. Thus, one is naturally led to models that preferentially alter $\mathcal{O}_9$ rather than $\mathcal{O}_{10}$.

Similarly, neither the radiative decay $B \to X_s\gamma$ nor the mass difference ΔM_s and mixing phase ϕ_s measurements for the B_s system show any appreciable discrepancy with the SM expectations. The pattern of deviations is thus a complicated one and, naively at least, does not appear to show a definite direction towards any well-motivated NP model. Consequently, most efforts at explaining the anomalies consider only a subset, either R_K and/or $R(D^{(*)})$ data [22,23] or $R_{K^{(*)}}$ and $b \to s\ell\ell$ data [24]. Those that do attempt a more complete treatment either invoke very complicated models or result in fits that are not very good. In addition, they are liable to result in other unacceptable phenomenological consequences. Analyses within specific models, like leptoquarks, are available in the literature [25].

In view of this, we adopt a very phenomenological approach rather than advocate a particular model. Assuming an effective Lagrangian, with the minimal number of new parameters, in the guise of the unknown Wilson coefficients (WCs), we seek the best fit. While not an entirely new idea, our analysis takes into account not only the anomalous channels but also the existing limits on several other channels; as we will show, they provide the tightest constraints on the parameter space. This approach hopefully will pave the way to unravelling the as yet unknown flavor dynamics.

Models.—Within the SM, the $b \to c\tau\bar{\nu}_\tau$ transition proceeds through a tree-level W exchange. If the NP adds coherently to the SM, one can write the effective Hamiltonian as

$$\mathcal{H}^{\rm eff} = \frac{4G_F}{\sqrt{2}} V_{cb}(1 + C^{\rm NP})[(c,b)(\tau,\nu_\tau)], \quad (6)$$

where the NP contribution is parametrized by $C^{\rm NP}$ vanishes in the SM limit and we have introduced the shorthand notation $(x, y) \equiv \bar{x}_L \gamma^\mu y_L \quad \forall~x, y$. To explain the data, one thus needs either small positive or large negative values of $C^{\rm NP}$.

The flavor-changing neutral-current decays $B \to K^{(*)}\mu\mu$ and $\phi\mu\mu$ are occasioned by the $b \to s\mu\mu$ transition proceeding, within the SM, primarily through a combination of the penguin and the box diagrams (driven, essentially, by the top quark). Parametrizing the ensuing effective Hamiltonian as

$$\mathcal{H}^{\rm eff} = \frac{-4G_F}{\sqrt{2}} V_{tb}V_{ts}^* \sum_i C_i(\mu)\mathcal{O}_i(\mu), \quad (7)$$

where the relevant operators are

$$\begin{aligned} \mathcal{O}_7 &= [\alpha_{\rm em}(m_b)m_b/4\pi](\bar{s}\sigma_{\mu\nu}P_R b)F^{\mu\nu}, \\ \mathcal{O}_9 &= [\alpha_{\rm em}(m_b)/4\pi](\bar{s}\gamma_\mu P_L b)(\bar{\mu}\gamma^\mu\mu), \\ \mathcal{O}_{10} &= [\alpha_{\rm em}(m_b)/4\pi](\bar{s}\gamma_\mu P_L b)(\bar{\mu}\gamma^\mu\gamma_5\mu). \end{aligned} \quad (8)$$

The WCs, matched with the full theory at m_W and then run down to m_b at the next-to-next-to-leading logarithmic accuracy [23], are given in the SM as $C_7 = -0.304$, $C_9 = 4.211$, and $C_{10} = -4.103$. The differential widths for the $B \to K^{(*)}\mu\mu$ decay are obtained in terms of algebraic functions of these. NP contributions to $\mathcal{H}^{\rm eff}$ can be parametrized by $C_i \to C_i + C_i^{\rm NP}$.

Similarly, the $b \to s\nu\bar{\nu}$ transition (which governs the $B \to K^{(*)}\nu\bar{\nu}$ decays) proceeds through the Z penguins and box diagrams. Unless right-handed neutrino fields are introduced, the low-energy effective Hamiltonian can be parametrized by [26]

$$\mathcal{H}^{\rm eff} = \frac{2G_F}{\sqrt{2}} V_{tb}V_{ts}^* \frac{\alpha_{\rm em}}{\pi} C_L^{\rm SM}(1 + C_\nu^{\rm NP})(s,b)(\nu,\nu), \quad (9)$$

where $C_\nu^{\rm NP}$ denotes the NP contribution. Including the NLO QCD correction and the two-loop electroweak contribution, the SM WC is given by $C_L^{\rm SM} = -X_t/s_w^2$, where the Inami-Lim function $X_t = 1.469 \pm 0.017$ [26,27].

While it may seem trivial to write down extra four-fermi operators that would produce just the right contributions, care must be taken to see that this does not introduce unwelcome consequences. For one, a large enhancement of C_{10} could lead to an unacceptably large $\mathrm{BR}(B_s \to \mu\mu)$, with $\mathcal{O}_{10}$ being the leading contributor to this decay. Similarly, the said four-fermi operators need to be invariant under the SM gauge group (assuming that the NP appears only above the electroweak scale). A nonzero $C^{\rm NP}$ [see Eq. (6)] would, potentially, lead to an analogue of $C_{10}^{\rm NP}$ for the tau channel. This, in turn, would lead to an enhancement of $B_s \to \tau\tau$, where the chirality suppression is less

operative than in the muonic case. Indeed, the LHCb Collaboration [28] has obtained a 95% C.L. upper limit of 6.8×10^{-3} on the branching fraction for this mode, with the SM value being $(7.73 \pm 0.49) \times 10^{-7}$ [20]. (It should be noted, though, that this analysis does not actually reconstruct the τ's but employs neural networks. Hence, it is possible that future measurements would point to a value higher than the limits quoted.) Similarly, none of the three operators $(b,s)(\nu_i,\nu_i)$ may receive large corrections lest the SM expectations, namely [26],

$$\mathrm{BR}(B^+ \to K^+ \nu \bar{\nu})_{\mathrm{SM}} = (3.98 \pm 0.43 \pm 0.19) \times 10^{-6},$$
$$\mathrm{BR}(B^0 \to K^{*0} \nu \bar{\nu})_{\mathrm{SM}} = (9.19 \pm 0.86 \pm 0.50) \times 10^{-6}, \quad (10)$$

be augmented (note that the neutrino flavors need not be identical for the NP) to levels beyond the 90% C.L. upper bounds (summed over all three neutrinos) as obtained by the Belle Collaboration [29], viz.

$$\mathrm{BR}(B \to K^{(*)} \nu \bar{\nu}) < 1.6(2.7) \times 10^{-5}. \quad (11)$$

In view of the aforementioned constraints, we consider only a combination of two four-fermi operators, characterized by a single WC (assumed to be real to avoid new sources of CP violation). Since we do not claim to obtain the ultraviolet completion thereof, we do not speculate on the (flavor) symmetry that would have led to such a structure, which could have arisen from a plethora of NP scenarios, such as models of (gauged) flavor, leptoquarks (or, within the supersymmetric paradigm, a breaking of R parity), etc. To wit, we propose a model involving two four-fermi operators in terms of the second- and third-generation (weak-eigenstate) fields

$$\mathcal{H}^{\mathrm{NP}} = A_1 (\bar{Q}_{2L} \gamma_\mu L_{3L})(\bar{L}_{3L} \gamma^\mu Q_{3L})$$
$$+ A_2 (\bar{Q}_{2L} \gamma_\mu Q_{3L})(\bar{\tau}_R \gamma^\mu \tau_R), \quad (12)$$

where the overall Clebsch-Gordan coefficients have been subsumed and we demand $A_2 = A_1$.

This operator, seemingly, contributes to $R(D^{(*)})$ but not to the other anomalous processes. This, though, is true only above the electroweak scale. Below this scale, the Hamiltonian needs to be rediagonalized. [With NP only modifying the Wilson coefficients of certain SM operators to a small extent, the QCD corrections (as well as hadronic uncertainties) are analogous. Additional effects due to operator mixings are too small to be of any concern.] In the quark sector, this is determined by the quark masses, and the small nonalignment due to $A_{1,2}$ can be neglected. In the leptonic sector, though, the extreme smallness of the neutrino masses implies that the nonuniversal term $\mathcal{H}^{\mathrm{NP}}$ plays a major role [30]. To this end, we consider the simplest of field rotations for the left-handed leptons from the unprimed (flavor) to the primed (mass) basis, namely,

$$\tau = \cos\theta\tau' + \sin\theta\mu', \qquad \nu_\tau = \cos\theta\nu_\tau' + \sin\theta\nu_\mu'. \quad (13)$$

This, immediately, generates a term with the potential to explain the $b \to s\mu\mu$ anomalies.

Results.—The scenario is, thus, characterized by two parameters, namely, A_1 and $\sin\theta$. The best-fit values for these can be obtained by effecting a χ^2 test defined through

$$\chi^2 = \sum_{i=1}^{7} \frac{(\mathcal{O}_i^{\mathrm{exp}} - \mathcal{O}_i^{\mathrm{th}})^2}{(\Delta\mathcal{O}_i^{\mathrm{exp}})^2 + (\Delta\mathcal{O}_i^{\mathrm{th}})^2}, \quad (14)$$

where $\mathcal{O}_i^{\mathrm{exp}}$ ($\mathcal{O}_i^{\mathrm{th}}$) denote the experimental (theoretical) mean and $\Delta\mathcal{O}_i^{\mathrm{exp}}$ ($\Delta\mathcal{O}_i^{\mathrm{th}}$) the corresponding 1σ uncertainty, with the theoretical values depending on the model parameters. We include a total of seven measurements for the evaluation of χ^2, namely, $R(D)$, $R(D^*)$, R_K, $R_{K^*}^{\mathrm{low}}$, $R_{K^*}^{\mathrm{cntr}}$, Φ, and $\mathrm{BR}(B_s \to \mu\mu)$ [while not affected by the NP interactions in Eq. (12), this is relevant for the scenario considered later]. Only for the last two observables do $\Delta\mathcal{O}_i^{\mathrm{th}}$ need to be considered explicitly, while, for the rest, they have been subsumed within the experimental results. For our numerical analysis, we use $V_{cb} = 0.0416$ and $V_{tb}V_{ts}^* = -0.0409$ and find, for the SM, $\chi^2_{\mathrm{SM}} \simeq 46$.

Within the new model, the best fit corresponds to $\chi^2_{\mathrm{min}} \simeq 9$ (denoting a marked improvement) with the NP contributions being $C_9^{\mathrm{NP}} = -1.7$ and $C^{\mathrm{NP}} = -2.12$. In terms of the model parameters, this corresponds to (note that there is a $\theta \to -\theta$ degeneracy)

$$A_1(= A_2) = -2.92 \text{ TeV}^{-2}, \qquad \sin\theta = \pm 0.022. \quad (15)$$

Even this low value of χ^2_{min} is largely dominated by a single measurement, namely, $R_{K^*}^{\mathrm{low}}$. This is not unexpected, as an agreement to this experimental value to better than 1σ is not possible if the NP contribution can be expressed just as a modification of the SM WCs rather than through the introduction of a new and small dynamical scale (such a change could be tuned so as to manifest itself primarily only in the low-q^2 region but is likely to have other ramifications). Note that the small value of $\sin\theta$ can only partially explain the atmospheric neutrino oscillation, while the full explanation needs additional dynamics.

More importantly, in effecting the field rotation of Eq. (13) in $\mathcal{H}^{\mathrm{NP}}$, we generate terms of the form $(s,b)(\mu,\tau)$, leading to potential lepton-flavor-violating (LFV) decays. The current limits on the relevant ones are [31]

$$\mathrm{BR}(B^+ \to K^+ \mu^\pm \tau^\mp) < 4.5(2.8) \times 10^{-5}. \quad (16)$$

In Fig. 1, we display the constraints from this particular mode. While the best-fit point is summarily ruled out, clearly solutions can be found if a slight worsening of the χ^2 (to $\simeq 15$) is acceptable. This would still represent a much better agreement than is possible within the SM. The corresponding values of the observables are $R_K = 0.86$,
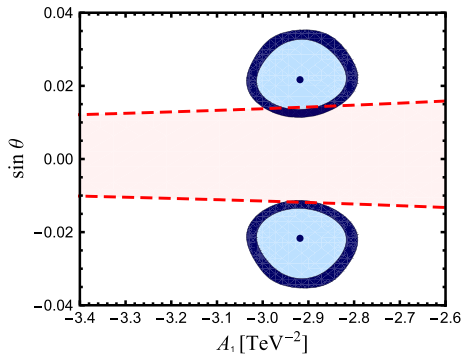
FIG. 1.   The light and dark blue regions denote 95% and 99% C.L. bands, respectively, around the best-fit points. The red shaded region is allowed by bounds from $\mathrm{BR}(B^+ \to K^+ \mu^- \tau^+)$.

$R_{K^*}^{\mathrm{cntr}} = 0.88$, $R_{K^*}^{\mathrm{low}} = 0.90$, $R(D^{(*)}) = 1.25 \times R_{\mathrm{SM}}(D^{(*)})$, and $\Phi = 4.1 \times 10^{-8}~\mathrm{GeV}^{-2}$, representing quite a reasonable fit to all but $R_{K^*}^{\mathrm{low}}$. It should be noted here that the $\theta \to -\theta$ degeneracy is broken by the LFV constraint, with $\theta > 0$ being slightly preferable.

Further improving the fit to $R_{K^{(*)}}$ requires the introduction of a small bit of C_{10}^{NP}. Postponing the discussion of $B_s \to \tau\tau$, this is most easily achieved if we choose to destroy, to a small degree, the relation $A_2 = A_1$. As an illustrative example, we consider $A_2 = 4A_1/5$. The consequent best-fit values for A_1 and $\sin\theta$ remain virtually the same but, now, $\chi^2_{\mathrm{min}} = 7$ with NP contributions being $C_9^{\mathrm{NP}} = -1.51$, $C_{10}^{\mathrm{NP}} = 0.17$, and $C^{\mathrm{NP}} = -2.12$. The result is depicted in Fig. 2. Once the LFV constraint is imposed, the observables at the overlap region are $R_K \simeq 0.80$, $R_{K^*}^{\mathrm{cntr}} \simeq 0.83$, $R_{K^*}^{\mathrm{low}} \simeq 0.88$, $R(D^{(*)}) \simeq 1.24 \times R_{\mathrm{SM}}(D^{(*)})$, and $\Phi \simeq 3.8 \times 10^{-8}~\mathrm{GeV}^{-2}$, showing marked improvement in the fit to all but $R_{K^*}^{\mathrm{low}}$ and correspond to $\chi^2 \simeq 10$. While the finite contribution to C_{10}^{NP} does enhance $B_s \to \tau\tau$, the latter (gray shaded region in Fig. 2) does not have a major impact. It should be realized, though, that a stronger breaking of the $A_2 = A_1$ relation would have led to a better (worse) agreement with the LFV ($B_s \to \tau\tau$) constraints.

It is interesting to speculate on the origin of this split between the A_i. A naive explanation would be to attribute the difference to the quantum numbers of the leptonic fields under an as yet unidentified gauge symmetry, with the attendant anomaly cancellation being effected by either invoking heavier fermionic fields or through other means. Care must be taken, however, not to induce undesirable phenomenology. An alternative is to attribute the difference to quantum corrections, although the aforementioned shift is somewhat larger than that expected from a naive renormalization group flow perspective, namely, $\sim (\alpha_{\mathrm{wk}}/4\pi) \ln(\Lambda_{\mathrm{NP}}^2/m_b^2)$, where $\Lambda_{\mathrm{NP}} \sim 1$ TeV is the putative scale of NP. It should be noted here, though, that the 20% shift is only illustrative and not really needed. Indeed, once the electroweak symmetry is broken, the various

pieces in $\mathcal{H}^{\mathrm{eff}}$ suffer differing renormalization group flow down to the m_b scale, and the consequent breaking of the degeneracy is, putatively, of the right magnitude to explain the remaining discrepancies.

It is worthwhile, at this stage, to explore the consequences of introducing other operators in $\mathcal{H}^{\mathrm{NP}}$. While operators constructed out of $SU(2)_L$-triplet currents (denoted by the subscript 3) such as $(\bar{Q}_{2L}\gamma^\mu Q_{3L})_3(\bar{L}_{3L}\gamma_\mu L_{3L})_3$, $(\bar{Q}_{2L}\gamma^\mu L_{3L})_3(\bar{L}_{3L}\gamma_\mu Q_{3L})_3$, etc., would also have admitted solutions to the anomalies, they, typically, would also result in unsuppressed $b \to s\nu\bar{\nu}$ transitions. Circumventing the bounds would, then, require the introduction of multiple operators and cancellations between them. We will discuss such possibilities in detail in a subsequent paper.

This would, typically, still leave behind too large a rate for $B_s \to \tau\tau$ (first reference of [25]) and, hence, needs the further introduction of yet another operator such as the second one in $\mathcal{H}^{\mathrm{NP}}$. Apart from enhancing $B_s \to \tau\tau$ ($B \to X_s\tau\tau$ and $\Lambda_b \to \Lambda\tau\tau$ are affected, too, but bounds from these sectors are not too serious), this would also affect the other modes to varying degrees. Consequently, the best-fit values will change. Indeed, a lower $\chi^2 (\simeq 5.4)$ is achievable for virtually the same A_1 but slightly smaller $|\sin\theta| (\simeq 0.018)$. Understandably, if both the $B_s \to \tau\tau$ bound as well that in Eq. (16) are to be satisfied, the χ^2 can be reduced to at most $\simeq 11$. Similarly, $\mathrm{BR}(B \to X_s\tau\tau)$, as well as $\mathrm{BR}(\Lambda_b \to \Lambda\tau\tau)$, will also be increased and should be close to observation at the LHCb. However, processes like $b \to s\gamma$ or $\tau \to \mu\gamma$ will remain under control, as we have checked. Similarly, while we do not "explain" $(g-2)_\mu$, the agreement is marginally better than within the SM. The new operators also generate, through renormalization group running, operators involving four leptons [32] and thus may lead to effects like $\tau \to 3\mu$. They are, however, well within control, mostly because of the small value of $\sin\theta$.

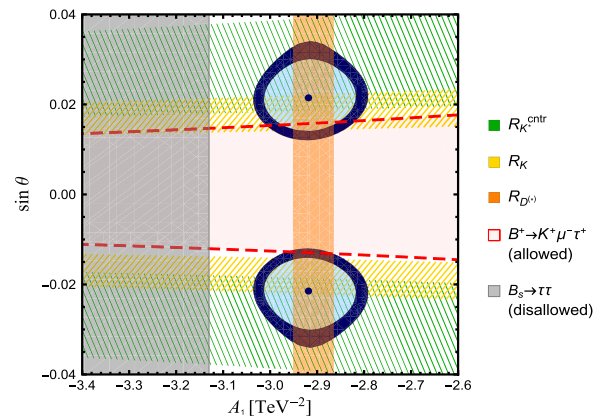FIG. 2.   The fit for $A_2 = 4A_1/5$, with the bands around the best-fit points corresponding to 95% and 99% C.L. Also shown are the 1σ bands from $R_{K^{(*)}}$ and $R(D)$ and the 95% upper limits from $B_s \to \tau\tau$ and $B^+ \to K^+\mu^-\tau^+$.

In summary, we have identified the minimal modification to the SM in terms of an effective theory that can explain the anomalies in both the charged- and the neutral-current decays of bottom mesons, a task that has been challenging on account of the seemingly contradictory requirements that the data demand. We circumvent this by postulating just two four-fermi operators with WCs related by a symmetry and taking advantage of the possibility of a small but nontrivial rotation of the charged lepton fields that a flavor-nonuniversal operator entails. Taking all the data into account, we find that with just two new parameters the χ^2 can be reduced from 46 (in the SM) to below 15 while being consistent with all other data. For the best-fit point, most observables are consistent within $\sim 1\sigma$, while R_{K^*} and BR($B \to \phi\mu\mu$) in the low-q^2 bins are consistent to only within $\sim 2\sigma$.

The scale of new physics that such an explanation demands is a few TeV at best, rendering searches at the LHC to be very interesting. An even stronger preference is that at least one of $B \to K^{(*)}\mu\tau$ and $B_s \to \tau\tau$ should be close to discovery. A more precise determination of the ratios that we have discussed in this Letter is, therefore, of prime importance, as this can open the door to new flavor dynamics and hence the world beyond the SM.

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