Flavor models for B[over ¯]→D^{(*)}τν[over ¯]

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Flavor models for $\bar{B} \to D^{(*)}\tau\bar{\nu}$

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The ratio of the measured $\bar{B} \to D^{(*)}\ell\bar{\nu}$ decay rates for $\ell = \tau$ vs. e, μ deviate from the Standard Model (SM) by about 4σ. We show that the data are in tension with the SM, independent of form factor calculations, and we update the SM prediction for $\mathcal{B}(B \to X_c\tau\bar{\nu})/\mathcal{B}(B \to X_c\ell\bar{\nu})$. We classify the operators that can accommodate the measured central values, as well as their UV completions. We identify models with leptoquark mediators that are minimally flavor violating in the quark sector, and are minimally flavor violating or τ-aligned in the lepton sector. We explore experimental signatures of these scenarios, which are observable in the future at ATLAS/CMS, LHCb, or Belle II.

I. INTRODUCTION

Measurements of the $\bar{B} \to D\tau\bar{\nu}$ and $\bar{B} \to D^*\tau\bar{\nu}$ decay rates are now available from BaBar [1, 2] and Belle [3] with their full datasets. The $\bar{B} \to D^*\tau\bar{\nu}$ decay mode was also observed recently by LHCb [4]. These measurements are consistent with each other and with earlier results [5, 6], and together show a significant deviation from Standard Model (SM) predictions for the combination of the ratios

$$R(X) = \frac{\mathcal{B}(\bar{B} \to X\tau\bar{\nu})}{\mathcal{B}(\bar{B} \to Xl\bar{\nu})}\,, \qquad (1)$$

where $l = e, \mu$. The measurements are consistent with e/μ universality [7, 8]. The $R(D^{(*)})$ data, their averages [9], and the SM expectations [10–12] are summarized in Table I. (If the likelihood of the measurements is Gaussian, then the deviation from the SM is more than 4σ.) Kinematic distributions, namely the dilepton invariant mass q^2, are also available from BaBar and Belle [2, 3], and must be accommodated by any model that modifies the rates. In the future, Belle II is expected to reduce the measured uncertainties of $R(D^{(*)})$ by factors of ~ 5 or more [13], thereby driving experimental and theory precision to comparable levels.

In the type-II two-Higgs-doublet model (2HDM), the $\bar{B} \to D^{(*)}\tau\bar{\nu}$ rate (as well as $B^- \to \tau\bar{\nu}$) receives contributions linear and quadratic in $m_b\, m_\tau \tan^2\beta/m_{H^\pm}^2$ [14–16], which can be substantial if $\tan\beta$ is large. However, the $R(D^{(*)})$ data are inconsistent with this scenario [1].

	$R(D)$	$R(D^*)$	Corr.
BaBar	$0.440 \pm 0.058 \pm 0.042$	$0.332 \pm 0.024 \pm 0.018$	-0.45
Belle	$0.375 \pm 0.064 \pm 0.026$	$0.293 \pm 0.038 \pm 0.015$	-0.32
LHCb		$0.336 \pm 0.027 \pm 0.030$	
Exp. average	0.388 ± 0.047	0.321 ± 0.021	-0.29
SM expectation	0.300 ± 0.010	0.252 ± 0.005	
Belle II, 50/ab	± 0.010	± 0.005	

TABLE I. Measurements of $R(D^{(*)})$ [1, 3, 4], their averages [9], the SM predictions [10–12], and future sensitivity [13]. The first (second) experimental errors are statistical (systematic).

Discovering new physics (NP) in transitions between the third and second generation fermion fields has long been considered plausible, since the flavor constraints are weaker on four-fermion operators mediating such transitions. (Prior studies of $B \to X_s\nu\bar{\nu}$ [17] and $B_{(s)} \to \tau^+\tau^-(X)$ [18, 19] decays were motivated by this consideration.) However, $\bar{B} \to D^{(*)}\tau\bar{\nu}$ is mediated by the tree-level $b \to c$ transition. It is suppressed in the SM neither by CKM angles (compared to other B decays) nor by loop factors, with only a modest phase space suppression due to the τ mass. This goes against the usual lore that the first manifestations of new physics at low energies are most likely to occur in processes suppressed in the SM.

The goal of this paper is to explore flavor structures for NP capable of accommodating the central values of the $R(D^{(*)})$ data summarized in Table I. To do so, a sizable NP contribution to semileptonic $b \to c$ decays must be present, and the NP mass scale must be near the weak scale. This requires nontrivial consistency with other constraints, such as direct searches at the LHC and precision electroweak data from LEP. When NP couplings to other generations are present, constraints from flavor physics, such as meson mixing and rare decays, also play a role. For example, any flavor model predicts some relation between the $\bar{b}c\,\bar{\nu}\tau$ and $\bar{b}u\,\bar{\nu}\tau$ operators, so models explaining $R(D^{(*)})$ must accommodate the observed $B^- \to \tau\bar{\nu}$ branching ratio, which agrees with the SM [20, 21]. We show below that despite strong constraints some scenarios remain viable and predict signals in upcoming experiments.

We begin by presenting new inclusive calculations that demonstrate that the measured central values of $R(D^{(*)})$ are in tension with the SM, independent of form factor computations. Then, in Sec. II, we perform a general operator analysis to identify which four-fermion operators simultaneously fit $R(D)$ and $R(D^*)$. In Sec. III we discuss possible mediators that can generate the viable operators. We identify working models with leptoquark mediators that are minimally flavor violating in the quark sector, and we confirm their consistency with current experimental constraints. Finally, Sec. IV contains our conclusions and a discussion of possible future signals at the LHC and Belle II. Appendix A contains a discussion of $U(2)^3$ models.

A. Standard Model considerations

The tension between the central values of the $R(D^{(*)})$ data and the SM is independent of the theoretical predictions for $R(D^{(*)})$ quoted in Table I. The measured $R(D^{(*)})$ values imply a significant enhancement of the inclusive $B \to X_c \tau\bar{\nu}$ rate, which can be calculated precisely in the SM using an operator product expansion, with theoretical uncertainties that are small and essentially independent from those of the exclusive rates.

To see this, note that the isospin-constrained fit for the branching ratios is quoted as [1]

$$\mathcal{B}(\bar{B} \to D^* \tau\bar{\nu}) + \mathcal{B}(\bar{B} \to D\tau\bar{\nu}) = (2.78 \pm 0.25)\% \,, \quad (2)$$

which applies for $B^\pm$ decays (recall the lifetime difference of $B^\pm$ and B^0). The averages in Table I imply for the same quantity the fully consistent result,

$$\mathcal{B}(\bar{B} \to D^* \tau\bar{\nu}) + \mathcal{B}(\bar{B} \to D\tau\bar{\nu}) = (2.71 \pm 0.18)\% \,. \quad (3)$$

The SM prediction for $R(X_c)$, the ratio for inclusive decay rates, can be computed in an operator product expansion. Updating results in Refs. [22, 23], and including the two-loop QCD correction [24], we find

$$R(X_c) = 0.223 \pm 0.004 \,. \quad (4)$$

The uncertainty mainly comes from m_b^{1S}, the HQET matrix element λ_1, and assigning an uncertainty equal to half of the order α_s^2 term in the perturbation series in the $1S$ scheme [25]. The most recent world average, $\mathcal{B}(B^- \to X_c e\bar{\nu}) = (10.92 \pm 0.16)\%$ [26, 27], then yields the SM prediction,

$$\mathcal{B}(B^- \to X_c \tau\bar{\nu}) = (2.42 \pm 0.05)\% \,. \quad (5)$$

In $B^- \to X_c e\bar{\nu}$ decay, hadronic final states other than D and D^* contribute about 3% to the 10.92% branching ratio quoted above, and the four lightest orbitally excited D meson states (often called collectively D^{**}) account for about 1.7%. Using Ref. [28] for the theoretical description of these decays, taking into account the phase space differences and varying the relevant Isgur-Wise functions, suggests $R(D^{**}) \gtrsim 0.15$ for the sum of these four states. This in turn implies for the sum of the central values of the rates to the six lightest charm meson states

$$\mathcal{B}(\bar{B} \to D^{(*)} \tau\bar{\nu}) + \mathcal{B}(\bar{B} \to D^{**}\tau\bar{\nu}) \sim 3\% \,, \quad (6)$$

in nearly 3σ tension with the inclusive calculation in Eq. (5). Note that Eqs. (2), (3), and (6) are also in mild tension with the LEP average of the rate of an admixture of b-flavored hadrons to decay to τ leptons [29],

$$\mathcal{B}(b \to X\tau^+\nu) = (2.41 \pm 0.23)\% \,. \quad (7)$$

Since both the experimental and theoretical uncertainties of $\mathcal{B}(B \to X_c \tau\bar{\nu})$ are different from the exclusive rates, its direct measurement from Belle and BaBar data would be interesting and timely [30].

II. $\bar{B} \to D^{(*)}\tau\bar{\nu}$ OPERATOR ANALYSIS

In this section we study operators mediating $b \to c\tau\bar{\nu}$ transitions. In contrast to prior operator fits [31–34], we adopt an overcomplete set of operators corresponding to all possible contractions of spinor indices and Lorentz structures to help with the classification of viable models. (We also take into account the constraints from q^2 spectra, which were unavailable at the time of the first operator analyses.) Although Fierz identities allow different spinor contractions to be written as linear combinations of operators with one preferred spinor ordering, the set of possible currents that can generate the operators is manifest in the overcomplete basis.

We parametrize the NP contributions by

$$\mathcal{H} = \frac{4G_F}{\sqrt{2}}\, V_{cb}\, \mathcal{O}_{V_L} + \frac{1}{\Lambda^2} \sum_i C_i^{(\prime,\prime\prime)}\, \mathcal{O}_i^{(\prime,\prime\prime)} \,. \quad (8)$$

(Throughout this paper we do not display Hermitian conjugates added to interaction terms as appropriate.) Here the primes denote different ways of contracting the spinors, as shown in Table II, which also presents their Fierz transformed equivalents in terms of the "canonically" ordered fields (unprimed operators). In the SM, only the $\mathcal{O}_{V_L}$ operator is present. (For illustration, the type-II 2HDM generates the operator $\mathcal{O}_{S_R}$ with $C_{S_R}/\Lambda^2 = -2\sqrt{2}\, G_F V_{cb}\, m_b\, m_\tau \tan^2\beta / m_{H^\pm}^2$.)

We do not consider the possibility of the neutrino being replaced by another neutral particle, such as a sterile neutrino, which yields additional operators. The large enhancement of an unsuppressed SM rate favors NP that can interfere with the SM. A non-SM field in the final state would preclude the possibility of interference, leading to larger Wilson coefficients and/or lower mass scales for the NP, making the interpretation in terms of concrete models more challenging.

We assume that the effects of NP can be described by higher dimension operators respecting the SM gauge symmetries. This is only evaded if the NP mediating these transitions is light or if it is strongly coupled at the electroweak scale; in either case there are severe constraints. We classify operators by the representations under $SU(3)_C \times SU(2)_L \times U(1)_Y$ of the mediators that are integrated out to generate them, as shown in the last column of Table II. Some mediators uniquely specify a

	Operator	Fierz identity	Allowed Current	$\delta\mathcal{L}_{\rm int}$
$\mathcal{O}_{V_L}$	$(\bar{c}\gamma_\mu P_L b)\,(\bar{\tau}\gamma^\mu P_L \nu)$		$(\mathbf{1},\mathbf{3})_0$	$(g_q \bar{q}_L \boldsymbol{\tau}\gamma^\mu q_L + g_\ell \bar{\ell}_L \boldsymbol{\tau}\gamma^\mu \ell_L) W'_\mu$
$\mathcal{O}_{V_R}$	$(\bar{c}\gamma_\mu P_R b)\,(\bar{\tau}\gamma^\mu P_L \nu)$			
$\mathcal{O}_{S_R}$	$(\bar{c} P_R b)\,(\bar{\tau} P_L \nu)$		$(\mathbf{1},\mathbf{2})_{1/2}$	$(\lambda_d \bar{q}_L d_R \phi + \lambda_u \bar{q}_L u_R i\tau_2 \phi^\dagger + \lambda_\ell \bar{\ell}_L e_R \phi)$
$\mathcal{O}_{S_L}$	$(\bar{c} P_L b)\,(\bar{\tau} P_L \nu)$			
$\mathcal{O}_T$	$(\bar{c}\sigma^{\mu\nu} P_L b)\,(\bar{\tau}\sigma_{\mu\nu} P_L \nu)$			
$\mathcal{O}'_{V_L}$	$(\bar{\tau}\gamma_\mu P_L b)\,(\bar{c}\gamma^\mu P_L \nu)$	$\longleftrightarrow \mathcal{O}_{V_L}$	$(\mathbf{3},\mathbf{3})_{2/3}$	$\lambda\,\bar{q}_L \boldsymbol{\tau}\gamma_\mu \ell_L \boldsymbol{U}^\mu$
$\mathcal{O}'_{V_R}$	$(\bar{\tau}\gamma_\mu P_R b)\,(\bar{c}\gamma^\mu P_L \nu)$	$\longleftrightarrow -2\mathcal{O}_{S_R}$	$(\mathbf{3},\mathbf{1})_{2/3}$	$(\lambda\,\bar{q}_L\gamma_\mu \ell_L + \tilde{\lambda}\,\bar{d}_R \gamma_\mu e_R) U^\mu$
$\mathcal{O}'_{S_R}$	$(\bar{\tau} P_R b)\,(\bar{c} P_L \nu)$	$\longleftrightarrow -\frac{1}{2}\mathcal{O}_{V_R}$		
$\mathcal{O}'_{S_L}$	$(\bar{\tau} P_L b)\,(\bar{c} P_L \nu)$	$\longleftrightarrow -\frac{1}{2}\mathcal{O}_{S_L} - \frac{1}{8}\mathcal{O}_T$	$(\mathbf{3},\mathbf{2})_{7/6}$	$(\lambda\,\bar{u}_R \ell_L + \tilde{\lambda}\,\bar{q}_L i\tau_2 e_R) R$
$\mathcal{O}'_T$	$(\bar{\tau}\sigma^{\mu\nu} P_L b)\,(\bar{c}\sigma_{\mu\nu} P_L \nu)$	$\longleftrightarrow -6\mathcal{O}_{S_L} + \frac{1}{2}\mathcal{O}_T$		
$\mathcal{O}''_{V_L}$	$(\bar{\tau}\gamma_\mu P_L c^c)\,(\bar{b}^c\gamma^\mu P_L \nu)$	$\longleftrightarrow -\mathcal{O}_{V_R}$		
$\mathcal{O}''_{V_R}$	$(\bar{\tau}\gamma_\mu P_R c^c)\,(\bar{b}^c\gamma^\mu P_L \nu)$	$\longleftrightarrow -2\mathcal{O}_{S_R}$	$(\bar{\mathbf{3}},\mathbf{2})_{5/3}$	$(\lambda\,\bar{d}_R^c \gamma_\mu \ell_L + \tilde{\lambda}\,\bar{q}_L^c \gamma_\mu e_R) V^\mu$
$\mathcal{O}''_{S_R}$	$(\bar{\tau} P_R c^c)\,(\bar{b}^c P_L \nu)$	$\longleftrightarrow \frac{1}{2}\mathcal{O}_{V_L}$	$(\bar{\mathbf{3}},\mathbf{3})_{1/3}$	$\lambda\,\bar{q}_L^c i\tau_2 \boldsymbol{\tau}\ell_L \boldsymbol{S}$
$\mathcal{O}''_{S_L}$	$(\bar{\tau} P_L c^c)\,(\bar{b}^c P_L \nu)$	$\longleftrightarrow -\frac{1}{2}\mathcal{O}_{S_L} + \frac{1}{8}\mathcal{O}_T$	$(\bar{\mathbf{3}},\mathbf{1})_{1/3}$	$(\lambda\,\bar{q}_L^c i\tau_2 \ell_L + \tilde{\lambda}\,\bar{u}_R^c e_R) S$
$\mathcal{O}''_T$	$(\bar{\tau}\sigma^{\mu\nu} P_L c^c)\,(\bar{b}^c\sigma_{\mu\nu} P_L \nu)$	$\longleftrightarrow -6\mathcal{O}_{S_L} - \frac{1}{2}\mathcal{O}_T$		

TABLE II. All possible four-fermion operators that can contribute to $\bar{B} \to D^{(*)}\tau\bar{\nu}$. Operators for which no quantum numbers are given can only arise from dimension-8 operators in a gauge invariant completion. For other operators the interaction terms which are subsequently integrated out are given. For the T operators we use the conventional definition of $\sigma^{\mu\nu} = i[\gamma^\mu, \gamma^\nu]/2$.

single operator, while others can generate two simultaneously. Anticipating the large Wilson coefficients necessary to fit the observed $R(D^{(*)})$ ratios, we focus on operators which can arise from dimension-6 gauge-invariant terms. Operators which can only come from SM gauge invariant dimension-8 terms or cannot be generated by integrating out a low-spin mediator will be omitted. Such contributions would be suppressed by additional powers of v/Λ, or could only arise from strongly coupled NP.

We calculate the contributions of all operators in the heavy quark limit [35]. Our method follows that of Ref. [36], and we rederived and confirmed those results. (A missing factor of $(1 - m_\ell^2/q^2)$ has to be inserted in Eq. (10) of Ref. [36].) We use the most precise single measurement of the h_{A_1} form factor [8], which equals the Isgur–Wise function in the heavy quark limit.

Higher order corrections are neglected, except for the following two effects that are known to be significant. For scalar operators a numerically sizable term, $(m_B + m_{D^*})/(m_b + m_c) \simeq 1.4$, arises from $\langle D^*|\bar{c}\gamma_5 b|B\rangle = -q^\mu \langle D^*|\bar{c}\gamma_\mu\gamma_5 b|B\rangle/(m_b + m_c)$. We also include the leading-log scale dependence of scalar (here C_S is either C_{S_L} or C_{S_R}) and tensor currents [37] in fits to models where they appear simultaneously,

$$C_S(m_b) = \left(\frac{\alpha_s(m_t)}{\alpha_s(m_b)}\right)^{-12/23} \left(\frac{\alpha_s(M)}{\alpha_s(m_t)}\right)^{-12/21} C_S(M)\,,$$
$$C_T(m_b) = \left(\frac{\alpha_s(m_t)}{\alpha_s(m_b)}\right)^{4/23} \left(\frac{\alpha_s(M)}{\alpha_s(m_t)}\right)^{4/21} C_T(M)\,. \quad (9)$$

For numerical calculations we use a reference scale $M =$ 750 GeV. The sensitivity to this choice is small, as most of the running occurs at low scales, between m_b and m_t.

To test the robustness of our results to $\mathcal{O}(\Lambda_{\rm QCD}/m_{c,b})$ corrections, we varied the slope parameter of the Isgur–Wise function, ρ^2, by ± 0.2 (motivated by Ref. [38]), and found less than 1σ change in the results. We leave consideration of $\mathcal{O}(\Lambda_{\rm QCD}/m_{c,b})$ corrections for the new physics for future work.

Figure 1 shows the results of χ^2 fits to $R(D)$ and $R(D^*)$ for each of the four-fermion operators in Table II individually. Here and below, our χ^2 includes experimental and SM theory uncertainties, but does not include theory uncertainties on NP, which are subdominant. Throughout, we assume that no new large sources of CP violation are present, i.e., we assume that the the phases of the NP operators are aligned with the phase of the SM vector operator. A contribution from the $\mathcal{O}_T$ operator or a modification of the SM contribution proportional to $\mathcal{O}_{V_L}$ (or any of its equivalents under Fierz identities) provide good fits to the data. The $\mathcal{O}_{S_L}$ operator can also fit the total rates, but it leads to q^2 spectra incompatible with observations. The operator $\mathcal{O}''_{S_L}$ also gives a good fit, which is not apparent from only considering the unprimed operators. (Note that the measurements of $R(D)$ and $R(D^*)$ depend on the operator coefficients, because the decay distributions are modified by the new physics contribution, affecting the experimental efficiencies and the measured rates [1, 39]. This effect cannot be included in our fits, providing another reason to take the χ^2 values as rough guides only.)

As noted earlier, certain mediators can generate two

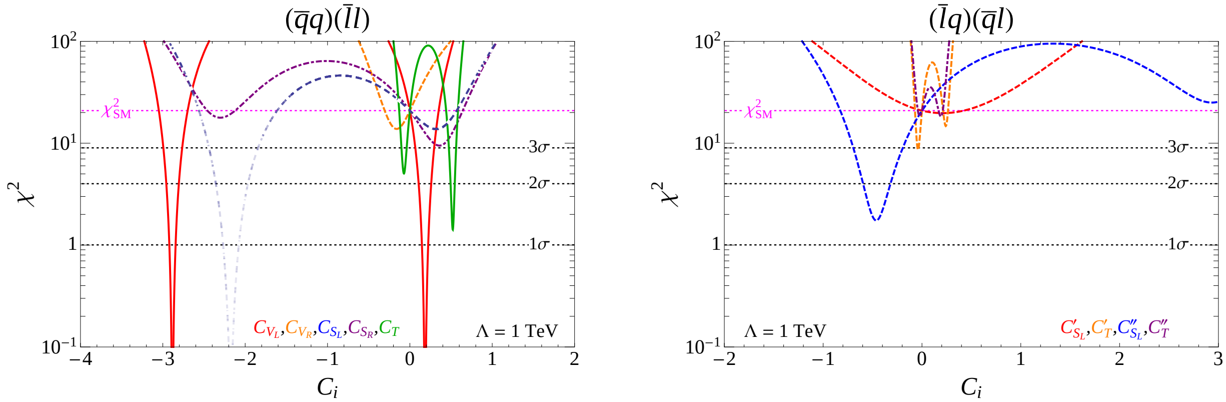

FIG. 1. Goodness-of-fit for the coefficients of individual operators from the measured $R(D)$ and $R(D^*)$ ratios. Besides the fits to the unprimed operators in Table II (left), we also show fits to primed operators not related by simple rescalings (right). Faded regions for C_{S_L} indicate good fits to the observed rates excluded by the measurement of the q^2 spectrum [2]. Note that the χ^2 includes experimental and SM theory uncertainties, but not theory uncertainties on NP.

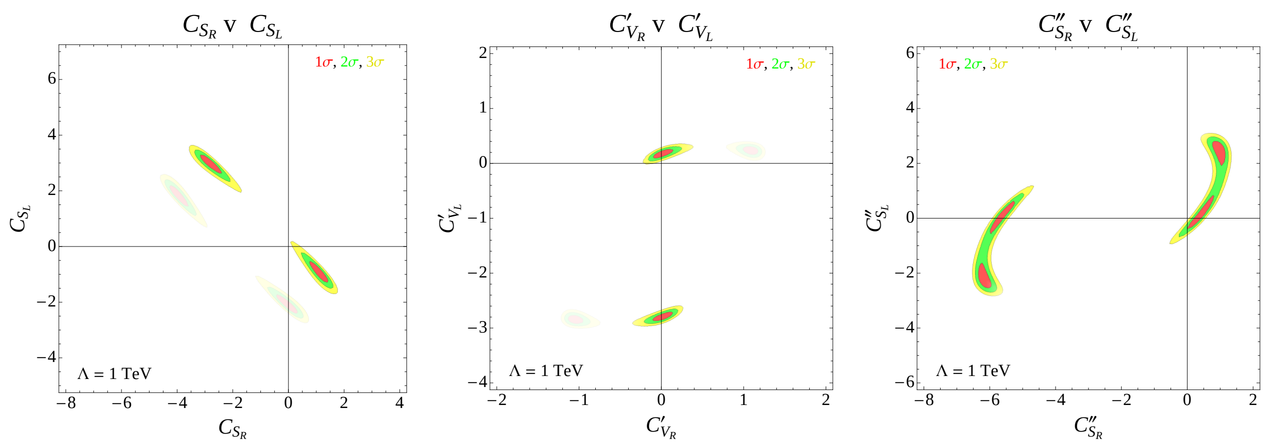

FIG. 2. Goodness-of-fit for coefficients of operators which can be generated from dimension-6 operators with fermion bilinears having the same SM quantum numbers. The plots show 1-, 2-, and 3σ allowed regions. Approximate regions of parameter space excluded by the measurement of the q^2 spectrum [2] are presented as faded regions, as in Fig. 1.

contributing operators simultaneously. Fig. 2 shows the three such two-dimensional χ^2 fits. While any two rates can be explained by fitting two operator coefficients, the existence of a solution consistent with all other constraints with a given flavor structure is nontrivial and is the topic of the following section. A summary of all coefficients of best fit points with $\chi^2_{\rm min} < 5$ and acceptable q^2 spectra is provided in Table III.

Besides the branching ratios, additional model discrimination comes from the q^2 spectra (especially in $\bar{B} \to D\tau\bar{\nu}$), which are consistent with SM expectations [2, 3]. It is not possible to do a combined fit with publicly available data, because correlations among different q^2 bins are unavailable. We follow Ref. [2] in eliminating certain models by comparing their predicted q^2 spectra with the measurement. It was observed that two of the four solutions in the C_{S_R}–C_{S_L} plane (Fig. 2, left plot) are excluded [2], as indicated by the faded regions. In the C'_{V_R}–C'_{V_L} plane (middle plot), we find the measured q^2 spectra exclude regions that provide good fits to the total rates for values of $|C'_{V_R}| \gtrsim 0.5$. In the C''_{S_R}–C''_{S_L} plane (right plot) all fits consistent with the total rates are also

Coefficient(s)	Best fit value(s) ($\Lambda = 1$ TeV)	
C_{V_L}	0.18 ± 0.04,	-2.88 ± 0.04
C_T	0.52 ± 0.02,	-0.07 ± 0.02
C''_{S_L}	-0.46 ± 0.09	
$(C_R,\, C_L)$	$(1.25, -1.02)$,	$(-2.84, 3.08)$
$(C'_{V_R},\, C'_{V_L})$	$(-0.01, 0.18)$,	$(0.01, -2.88)$
$(C''_{S_R},\, C''_{S_L})$	$(0.35, -0.03)$,	$(0.96, 2.41)$,
	$(-5.74, 0.03)$,	$(-6.34, -2.39)$

TABLE III. Best-fit operator coefficients with acceptable q^2 spectra and $\chi^2_{\rm min} < 5$. For the 1D fits in Fig. 1 we include the $\Delta\chi^2 < 1$ ranges (upper part), and show the central values of the 2D fits in Fig. 2 (lower part).

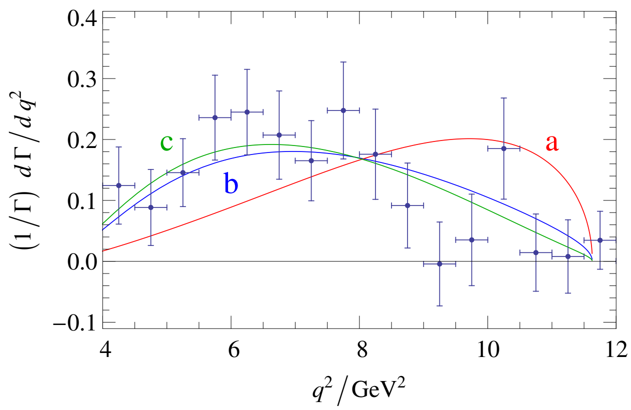

FIG. 3. Normalized $d\Gamma(\bar{B} \to D\tau\bar{\nu})/dq^2$ distributions. The histogram shows the BaBar data [2]. The red, blue, and green curves correspond to models with Wilson coefficients in the (a) top, (b) middle, and (c) bottom lines in Table IV.

Label	Coefficients ($\Lambda = 1$ TeV)		Comment
(a)	$C'_{V_R} = 1.10$	$C'_{V_L} = 0.24$	disfavored
(b)	$C'_{V_R} = -0.01$	$C'_{V_L} = 0.18$	allowed
(c)	$C''_{S_R} = 0.96$	$C''_{S_L} = 2.41$	allowed

TABLE IV. Operator coefficients for the q^2 spectra in Fig. 3.

consistent with the q^2 spectra. Figure 3 illustrates this by showing the measured $d\Gamma(B \to D\tau\bar{\nu})/dq^2$ spectrum together with one disfavored and two viable models, corresponding to the entries in Table IV.

III. MODELS

Having identified dimension-6 operators that can generate the observed $\bar{B} \to D^{(*)}\tau\bar{\nu}$ rate, we now consider possible UV completions. The TeV-scale required for the four-fermion operators points to a tree-level NP contribution. Concerning the flavor structure, one possibility is that NP is aligned with the SM Yukawa matrices and in the fermion mass basis only gives rise to the $\bar{b}c\,\bar{\nu}\tau$ four-fermion interaction. In such a scenario the measurement of $R(D^{(*)})$ would have little interplay with other flavor data; however, it is hard to imagine a UV completion with such precise alignment with the SM flavor structure. On the other hand, if operators are generated with general fermion content, $P_{ijkl}\,\bar{d}_i u_j\,\bar{\nu}_k \ell_l$, where i, j, k, l are generation indices, then some P_{ijkl} coefficients must satisfy strong experimental constraints. Most of the completions we consider below have also been studied, but without considerations to their possible flavor structure; see, e.g., Refs. [31–33, 40–44]

Independent of the UV completion, flavor-anarchic couplings are excluded by other rare decays and CKM unitarity constraints [45]. We therefore consider NP that is charged under a flavor symmetry, focusing on the possibility of Minimal Flavor Violation (MFV) [46–48]. In the quark sector, the MFV framework assumes that the breaking of the $U(3)_Q \times U(3)_u \times U(3)_d$ flavor symmetry has a single source in both the SM and beyond. It is parametrized by the Yukawas, acting as spurions of the symmetry breaking, transforming in the $(\mathbf{3}, \bar{\mathbf{3}}, \mathbf{1})$ representation for Y_u and the $(\mathbf{3}, \mathbf{1}, \bar{\mathbf{3}})$ for Y_d. Focusing primarily on MFV in the quark sector, for viable models we attempt to extend the MFV scenario to the lepton sector as well. We find a unique model that is consistent with MFV both in the quark and the lepton sector, and comment on alternative approaches involving horizontal symmetries. We begin by showing in Sec. III A that uncolored mediators are disfavored by other constraints. Then, in Sec. III B, we identify viable MFV models with leptoquark mediators.

A. Uncolored mediators

1. *Higgs-like scalars*

We first explain why no new color-neutral scalar, such as those in the nonstandard 2HDMs proposed in Refs. [49–51], is compatible with an MFV structure. The simplest case, a flavor singlet scalar, would be constrained to have couplings proportional to the SM Higgs. As shown in Fig. 2, comparable values of C_{S_L} and C_{S_R} are needed to explain the current $\bar{B} \to D^{(*)}\tau\bar{\nu}$ data. C_{S_L} is proportional to y_c, which implies that $\bar{B} \to D^{(*)}\tau\bar{\nu}$ cannot be fit, keeping the charged Higgs heavier than collider limits, consistently with perturbativity of y_t.

An alternative is to charge the scalar under the quark flavor symmetries. In this case, the scalar needs to be in the representation of some combination of Yukawa spurions, since it must couple to leptons as well. To generate both scalar couplings simultaneously, an unsuppressed $\mathcal{O}(1)$ coupling to one chiral combination of first-generation quarks is unavoidable. Integrating out the scalar generates dangerous four-quark and two-quark two-lepton operators. Four-quark operators of this type are strongly constrained, and such a scalar would also appear in $\tau^+\tau^-$ resonance searches [52, 53].

2. *W′-like vector triplets*

A rescaling of the SM operator, $\mathcal{O}_{V_L}$, provides a good fit to the data. A simple possibility is then the presence of a W', a new vector that couples similarly to the SM W boson. To avoid explicitly breaking the SM gauge symmetry, a full $SU(2)_L$ triplet is needed. The simplest choice, a flavor singlet W', is tightly constrained by the LHC [54–56], with electroweak-strength coupling of a W' to first generation quarks excluded up to $m_{W'} \sim 1.8$ TeV. This is in conflict with the $0.2 \sim g^2\,|V_{cb}|\,(1~\text{TeV}/m_{W'})^2$ coupling needed to explain the $\bar{B} \to D^{(*)}\tau\bar{\nu}$ data.

As with the scalars, it is possible to consider a W' that transforms under a nontrivial representation of the $[U(3)]^3$ quark flavor symmetry group, in an attempt to suppress the couplings to light quark generations. However, the W' needs to simultaneously couple to leptons, fixing its flavor assignment to be the conjugate of some combination of Yukawa spurions. This means that the same flavor structure allowing the W' to couple to leptons can always couple to a current of left-handed quarks contracted as a flavor singlet, allowing for universal couplings to all quark generations.

Suppose that the flavor singlet universal coupling to quarks is accidentally small. The minimal charge assignments are $(\bar{\mathbf{3}},\mathbf{3},\mathbf{1})$ and $(\bar{\mathbf{3}},\mathbf{1},\mathbf{3})$. In the $(\bar{\mathbf{3}},\mathbf{3},\mathbf{1})$ case, (semi)leptonic $b \to c$ transitions are suppressed by y_c, and therefore the central values of $R(D^{(*)})$ cannot be generated with perturbative couplings. In the $(\bar{\mathbf{3}},\mathbf{1},\mathbf{3})$ case, the $R(D^{(*)})$ data can be accommodated by the interaction $\delta\mathcal{L} = g\,\bar{q}_L^i Y_d^{ij} \boldsymbol{\tau}\cdot \boldsymbol{W}_\mu^{jk}\gamma^\mu q_L^k$ and an equivalent coupling to the leptons with the flavor indices of Y_d and $\boldsymbol{W}_\mu$ contracted into each other, given $y_b = \mathcal{O}(1)$. LHC direct production bounds are then evaded. However, tree-level FCNCs are mediated by the neutral component, W^0, and fitting the $\bar{B} \to D^{(*)}\tau\bar{\nu}$ rates implies too large of a contribution to B–$\bar{B}$ mixing. For an approach using a combination of dynamical singlet suppression and reduced symmetries to evade these issues, see Ref. [57].

B. Theories of MFV leptoquarks

The second class of potential tree-level mediators are leptoquarks. With no way to contract a single quark representation with Yukawa spurions to form a flavor singlet, leptoquarks in the MFV framework must carry some charges under the quark flavor symmetries for flavor indices to be contracted in interaction terms. Additionally, for representations of the flavor group to be coupled together in a way that gives non-zero couplings, the mediators must come in fundamental representations of some of the $[U(3)]^3$ quark flavor group. Any larger representation will always either give vanishing coupling when contracted with Yukawa spurions in order to couple to quarks or have the spurions entirely contracted into the leptoquark so as to not generate any new couplings. For reviews of leptoquarks, see, e.g., Refs. [58, 59], and for discussion of MFV leptoquarks, see, e.g., Ref. [60].

The six simplest possible flavor assignments under $U(3)_Q \times U(3)_u \times U(3)_d$, for scalar leptoquarks, are

$$S \sim (\bar{\mathbf{3}},\mathbf{1},\mathbf{1})\,, \quad S \sim (\mathbf{1},\bar{\mathbf{3}},\mathbf{1})\,, \quad S \sim (\mathbf{1},\mathbf{1},\bar{\mathbf{3}})\,, \tag{10}$$

while vector leptoquarks can have charge

$$U_\mu \sim (\mathbf{3},\mathbf{1},\mathbf{1})\,, \quad U_\mu \sim (\mathbf{1},\mathbf{3},\mathbf{1})\,, \quad U_\mu \sim (\mathbf{1},\mathbf{1},\mathbf{3})\,. \tag{11}$$

For each of these six charge assignments, the leptoquarks could be either electroweak $SU(2)_L$ singlets or triplets. Unless otherwise specified, all comments below apply for both cases. Since these leptoquarks need to mediate processes involving different generations of quarks, additional experimental constraints are important. We will see that the dominant bounds come from the tension between allowing for a large enough effect to fit the $R(D^{(*)})$ data, while being consistent with precision electroweak constraints and contributions to FCNC decays, such as $\bar{B} \to X_s\nu\bar{\nu}$.

Most choices of representation under the flavor groups can be immediately discarded, since they either only generate Yukawa-suppressed operators for the required $b \to c\tau\bar{\nu}$ decays, or have unsuppressed couplings to first-generation quarks. The latter case would lead to new contributions to $\tau^+\tau^-$ production at the LHC, which is severely constrained by Z' searches [52, 53]. We estimate that couplings larger than $\lambda/m \sim 0.25/\text{TeV}$ are excluded due to production of $\tau^+\tau^-$ through t-channel leptoquark exchange. Below, we only consider models that satisfy both of these constraints.

1. Scalar leptoquarks — τ alignment

Consider a scalar leptoquark, S, with $U(3)_Q \times U(3)_u \times U(3)_d$ charge $(\mathbf{1},\mathbf{1},\bar{\mathbf{3}})$. The interaction terms with the smallest number of Yukawa insertions for the electroweak singlet case take the form

$$\begin{aligned}\delta\mathcal{L} &= S(\lambda Y_d^\dagger\, \bar{q}_L^c i\tau_2 \ell_L + \tilde{\lambda} Y_d^\dagger Y_u\, \bar{u}_R^c e_R) \\ &= S_i(\lambda y_{d_i} V_{ji}^*\, \bar{u}_{Lj}^c \tau_L - \lambda y_{d_i} \bar{d}_{Li}^c \nu_L + \tilde{\lambda} y_{d_i} y_{u_j} V_{ji}^*\, \bar{u}_{Rj}^c \tau_R)\,.\end{aligned} \tag{12}$$

Couplings to light quarks are suppressed by quark masses and/or small CKM angles, so LHC bounds involving production off valence quarks are evaded. Meanwhile, $b \to c$ operators receive contributions of the form

$$\frac{C''_{S_R}}{\Lambda^2} = \frac{V_{cb}^*}{m_{S_3}^2}\lambda^2 y_b^2\,, \qquad \frac{C''_{S_L}}{\Lambda^2} = \frac{V_{cb}^*}{m_{S_3}^2}\lambda\tilde{\lambda}\, y_c y_b^2\,. \tag{13}$$

This contribution can be sizable if $y_b = \mathcal{O}(1)$, which can follow from a type II 2HDM at large $\tan\beta$. Due to the large suppression of C''_{S_L} from y_c, we use from Fig. 2, $C''_{S_L} \approx 0$ and $C''_{S_R}/\Lambda^2 = (0.35 \pm 0.1)/(1\ \text{TeV})^2$, corresponding to $\lambda y_b/m_{S_3} \approx (2.95 \pm 0.6)/\text{TeV}$, with $\tilde{\lambda}$ arbitrary as long as it remains perturbative. The analogous term for the electroweak triplet scalar only generates the C''_{S_R} term, but with the opposite sign, due to the presence of an extra τ^3 factor in the interaction Lagrangian for the relevant component. We do not study the triplet further, since fitting the $R(D^{(*)})$ data would require a large imaginary part for λ. (In this paper we neglect the possibility of complex couplings; otherwise CP violation constraints need to be studied, which go beyond the scope of this paper and are unlikely to result in qualitatively different allowed scenarios.) Similar operators contribute to $B^- \to \tau\bar{\nu}$ decay, with V_{cb} replaced by V_{ub}. The contribution of these operators to the measured rate is consistent with current data.

Coupling exclusively to third generation leptons can be enforced by imposing a horizontal flavor symmetry, $U(1)_\tau$, under which the leptoquark and third-generation lepton fields are equally and oppositely charged. Neutrino masses and mixings require additional spurions that break the lepton flavor symmetry, but these effects can be small enough to ignore in the present discussion. For an example where such symmetries generate the entire structure of $b \to c\tau\bar\nu$ transitions, see Ref. [61], and Refs. [62, 63] for discussions of lepton symmetries in $b \to c$ transitions in relation to other anomalies.

The presence of a large $S_3\bar{t}^c\tau$ coupling also generates a deviation from the SM prediction for the $Z\tau^+\tau^-$ coupling, which can affect both the total rate and asymmetry of $Z \to \tau^+\tau^-$ decays. For scalar leptoquarks, such corrections are finite and calculable [64]. In the case where only the couplings proportional to λ are present, these provide a bound on $\lambda y_b/m_{S_3}$ comparable to the value needed to fit the $R(D^{(*)})$ data. However, interference between the λ and $\tilde\lambda$ couplings can be large, since the latter is not suppressed by y_c (unlike in the decay of the b-quark). If $|\tilde\lambda/\lambda| \gtrsim 0.45$, interference among λ and $\tilde\lambda$ contributions can render the model unobservable to current electroweak precision tests.

Being color triplets, leptoquarks can be pair-produced directly at hadron colliders from gg initial states, independent of their couplings to fermions. The leptoquark that contributes to $\bar{B} \to D^{(*)}\tau\nu$, S_3, decays almost entirely to either $t\tau$ or $b\nu$ final states, with relative branching fraction determined by the λ and $\tilde\lambda$ couplings, but with a branching fraction to $t\tau$ not less than ≈ 0.5. Such searches have been carried out in both channels by CMS [65]. Assuming equal branching fractions for a conservative bound, the mass of the S_3 leptoquark is required to be $m_{S_3} \gtrsim 560$ GeV. Limits from decay purely to the $b\nu$ final state can also be derived by reinterpreting the bounds from sbottom searches [66, 67], providing comparable limits in areas of parameter space where they are applicable.

While the leading MFV interactions do not mediate a contribution to $b \to s\nu\bar\nu$, the SM rate is sufficiently small that we also need to consider effects subleading in the MFV expansion. The interactions in Eq. (12) refer to the mass eigenstates only if the leptoquarks' couplings to different down-type quarks do not mix. If such mixing occurs, it could induce large $b \to s\nu\bar\nu$ transition. In a spurion analysis, we should consider any number of spurion insertions. Since we need y_b to be $\mathcal{O}(1)$, we should include spurions of Y_d at all orders [68], and the above conclusions should be checked after allowing such insertions. Thus, the leptoquark mass matrix takes the form (again in the quark mass basis)

$$\begin{aligned} m_S^2 &= m_0^2\,(I + bY_d^\dagger Y_d + \cdots) \\ \Rightarrow m_{S_i}^2 &= m_0^2\,(1 + by_{d_i}^2 + \cdots)\,. \end{aligned} \tag{14}$$

This makes it clear that the leptoquarks do not mix to all orders in the MFV expansion, while the third generation leptoquark can be split in mass from the first two by a sizable amount.

Although the leptoquark mass matrix does not induce large $b \to s\nu\bar\nu$ transitions, strong constraints arise from interactions at next order in the MFV expansion.[1] At leading order, Eq. (12) couples down-type quarks to neutrinos, but as explained above, does not allow them to mix. The next term, with three Yukawa spurion insertions,

$$\delta\mathcal{L}' = \lambda' S Y_d^\dagger Y_u Y_u^\dagger\, \bar{q}_L^c i\tau_2 \ell_L\,, \tag{15}$$

generates interactions of the form

$$\delta\mathcal{L}' = S_i\,\lambda'\,y_{d_i} V_{ji}^*\, y_{u_j}^2\,\left(\bar{u}_{Lj}^c\tau_L - V_{jk}\,\bar{d}_{Lk}^c\nu_L\right). \tag{16}$$

This interaction generates $b \to s\nu\bar\nu$ decay, to which the leading contribution is

$$\mathcal{O}_{bs\nu\bar\nu} = -\frac{y_t^2 y_b^2\,\lambda\lambda'}{2m_{S_3}^2}\,V_{tb}^* V_{ts}\,\left(\bar{b}_L\gamma^\mu s_L\,\bar\nu_L\gamma_\mu\nu_L\right)\,, \tag{17}$$

with no parametric suppression relative to $b \to c$. Such transitions are most strongly constrained at present by [69]

$$\mathcal{B}(B^+ \to K^+\nu\bar\nu) < 1.6\times 10^{-5}\,. \tag{18}$$

Comparing the SM prediction with this bound [70, 71] gives $\lambda'\lambda y_b^2/(2m_{S_3}^2) \lesssim 0.26/(1\text{ TeV})^2$. If C''_{S_R} in Eq. (13) is to explain the values of $R(D^{(*)})$, this can be rewritten as $\lambda'/\lambda \lesssim 0.1$. Thus, a moderate suppression of the subleading coefficient of the MFV expansion is necessary to avoid the $b \to s\nu\bar\nu$ constraint.

2. *Scalar leptoquarks — lepton MFV*

We can extend the MFV framework to the lepton sector, and give the leptoquark an appropriate charge under the $U(3)_L \times U(3)_e$ symmetry broken by the lepton Yukawa couplings. Assigning the leptoquark to either of the lepton flavor representations $(\bar{\mathbf{3}},\mathbf{1})$ or $(\mathbf{1},\bar{\mathbf{3}})$ ensures that either the terms with λ or $\tilde\lambda$, respectively, will couple to all lepton generations uniformly. In the first case, the leptoquark contributes as much to $B \to D^{(*)}l\bar\nu$ as to final states with a τ, up to effects of subleading spurion insertions, and $R(D^{(*)})$ cannot deviate significantly from the SM prediction. In the second case, all couplings generating corrections to $\bar{B} \to D^{(*)}l\bar\nu$ are suppressed by Yukawas of either lighter quarks or leptons.

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[1] Due to the large value of y_t, terms to all orders in $(Y_uY_u^\dagger)^n$ should also be included in the expansion analogously to the $(Y_dY_d^\dagger)^n$ terms in Eq. (14). Up to terms suppressed by m_c/m_t, however, this only leads to a rescaling of the coefficient in the first subleading term.

For this second case, the leading interaction terms take the form

$$\begin{aligned}\delta\mathcal{L} &= S_{il}(\lambda Y^*_{d\,ji} Y_{eml}\, \bar{q}^c_{Lj} i\tau_2 \ell_{Lm} + \tilde{\lambda} Y^*_{d\,ji} Y_{ujk}\, \bar{u}^c_{Rk} e_{Rl}) \\ &= S_{il}(\lambda y_{d_i} y_{e_l} V^*_{ji}\, \bar{u}^c_{Lj}\tau_{Ll} - \lambda y_{d_i} y_{e_l} \bar{d}^c_{Li}\nu_{Ll} \\ &\qquad + \tilde{\lambda} y_{d_i} y_{u_j} V^*_{ji}\, \bar{u}^c_{Rj}\tau_{Rl})\,. \end{aligned} \tag{19}$$

The leptoquarks coupling to lighter lepton generations introduce additional constraints. Of those involving the λ coupling, the leading constraint comes from considering $D^0 \to \mu^+\mu^-$ decays, which receive a tree-level contribution from t-channel leptoquark exchange. However, the contribution to the branching ratio is suppressed by $\mathcal{O}(y_\mu^2)$ (in addition to the already-present helicity suppression), making it two orders of magnitude below current bounds [72] when consistent with the fit to $R(D^{(*)})$. A contribution to $B_s \to \mu^+\mu^-$ is generated by a mixed W–S_{32} box diagram, but the resulting correction to the SM rate is at the percent level. A more severe constraint on $\tilde{\lambda}$ exists from corrections to R_μ, measured at LEP, due to the unsuppressed coupling of the $S\bar{t}^c_R\ell_R$ vertex, providing a limit $\tilde{\lambda}/\lambda \lesssim 0.8$.

The leptoquarks' coupling to light leptons mean that collider searches for first- and second-generation leptoquarks become relevant. These bounds are less constraining then the third-generation searches in Sec. III B 1, however, since they involve different leptoquark components than that contributing to $R(D^{(*)})$. Here, the mass matrix takes the form

$$\begin{aligned}(m_S^2)_{ijlm} &= m_0^2\,(I + bY^*_{d\,ki}Y_{dkj} + \cdots + dY^*_{e\,kl}Y_{ekm} + \cdots) \\ \Rightarrow m^2_{S_{il}} &= m_0^2\,(1 + by^2_{d_i} + dy^2_{e_l} + \cdots)\,, \end{aligned} \tag{20}$$

so the S_{33} component can be lighter than the others by an $\mathcal{O}(1)$ amount without a sizable tuning. The fact that the bound on doubly-produced first- and second-generation leptoquarks are approximately a factor of 2 more stringent than those for the third-generation ones [73, 74] does not then lead to any unavoidable additional constraints. Furthermore, the additional leptoquarks do not lead to more stringent constraints from $b \to s\nu\bar{\nu}$ decays, since all couplings to neutrinos other than ν_τ are Yukawa suppressed.

3. *Vector leptoquarks*

Consider a vector leptoquark, U^μ, with $U(3)_Q \times U(3)_u \times U(3)_d$ charge $(\mathbf{1},\mathbf{1},\mathbf{3})$. If the leptoquark is an electroweak singlet, the lowest order MFV contribution in the quark mass basis is

$$\begin{aligned}\delta\mathcal{L} &= (\lambda\,\bar{q}_L Y_d \gamma_\mu \ell_L + \tilde{\lambda}\,\bar{d}_R\gamma_\mu e_R)\,U^\mu \\ &= \left(\lambda y_{d_i} V_{ji}\,\bar{u}_{Lj}\gamma_\mu\nu_L + \lambda y_{d_i}\bar{d}_{Li}\gamma_\mu\tau_L + \tilde{\lambda}\bar{d}_{Ri}\gamma_\mu\tau_R\right) U_i^\mu\,, \end{aligned} \tag{21}$$

while for the electroweak triplet, it is

$$\begin{aligned}\delta\mathcal{L} &= (\lambda\,\bar{q}_L Y_d \vec{\tau}\gamma_\mu\ell_L)\,\vec{U}^\mu \\ &= \left(\lambda y_{d_i} V_{ji}\,\bar{u}_{Lj}\gamma_\mu\nu_L + \lambda y_{d_i}\bar{d}_{Li}\gamma_\mu\tau_L\right) U_i^{0\mu} \\ &\quad + \frac{\lambda}{\sqrt{2}}\left(y_{d_i}\bar{d}_{Li}\gamma_\mu\nu_L\, U_i^{+\mu} + y_{d_i}V_{ji}\,\bar{u}_{Lj}\gamma_\mu\tau_L\, U_i^{-\mu}\right). \end{aligned} \tag{22}$$

(Here the signs on U do not indicate charges, but rather which component of the triplet we refer to.) The contribution to $b \to c$ decays comes from

$$\frac{C'_{V_L}}{\Lambda^2} = \frac{V_{cb}}{m^2_{U_3}}\,\lambda^2 y_b^2\,. \tag{23}$$

As in the prior case, this leptoquark model requires $y_b = \mathcal{O}(1)$ to be able to explain the $R(D^{(*)})$ data with perturbative couplings. In this case, it is the electroweak singlet that generates an operator with the wrong sign to be fit with real values of λ. Using Table II and Fig. 1, we see that a value of $\lambda y_b/m_{U_3} \approx (2.2 \pm 0.4)/\mathrm{TeV}$ gives a good fit.

If we consider giving leptonic flavor charges to the leptoquarks, as in the scalar case, then the $(\bar{\mathbf{3}},\mathbf{1})$ representation prevents a deviation of $R(D^{(*)})$ from being generated as before. The $(\mathbf{1},\bar{\mathbf{3}})$ representation does not, but in this case, a tree-level contribution to $B_s \to \mu^+\mu^-$ [75, 76] is generated at leading order in the MFV expansion. We estimate this to give a constraint of $\lambda y_b/m_{U_3} \lesssim 1.0/\mathrm{TeV}$, in significant tension with the value required to fit the B decay data. However, coupling to purely τ can be enforced via a horizontal symmetry, as in Sec. III B 1, and this is what we assume for the rest of the discussion.

As in the scalar case, we need to consider the possibility of generating sizable contributions to $b \to s\nu\bar{\nu}$ from additional spurion insertions. The leptoquark mass matrix remains diagonal. As in the scalar case, for the electroweak triplet vector leptoquark, the next contribution to the MFV expansion of the interaction term, with three Yukawa spurion insertions,

$$\delta\mathcal{L}' = \lambda'\,\bar{q}_L Y_u Y_u^\dagger Y_d\,\vec{\tau}\gamma_\mu\ell_L\vec{U}^\mu, \tag{24}$$

leads to an operator of the form

$$\mathcal{O}_{bs\nu\bar{\nu}} = \frac{V^*_{tb}V_{ts}}{m^2_{U_3}}\,y_t^2 y_b^2\,\lambda'\lambda\,\left(\bar{b}_L\gamma^\mu s_L\,\bar{\nu}_L\gamma_\mu\nu_L\right). \tag{25}$$

In this case, the bound is $\lambda'\lambda y_b^2/m^2_{U_3} \lesssim 0.13/(1\ \mathrm{TeV})^2$. If the operators in Eq. (23) are to explain the $R(D^{(*)})$ data, this can be rewritten as $\lambda'/\lambda \lesssim 0.05$. The electroweak triplet case thus requires a suppression of the subleading coefficient of the MFV expansion to avoid violating the $b \to s\nu\bar{\nu}$ bound. There is no such constraint for the electroweak singlet vector leptoquark, however, that model requires complex couplings to fit the $R(D^{(*)})$ data.

There are bounds on the leptoquarks produced directly through their gauge couplings via $gg \to UU$ pair production. The leptoquarks can then decay to $b\tau$ or $t\nu$ pairs.

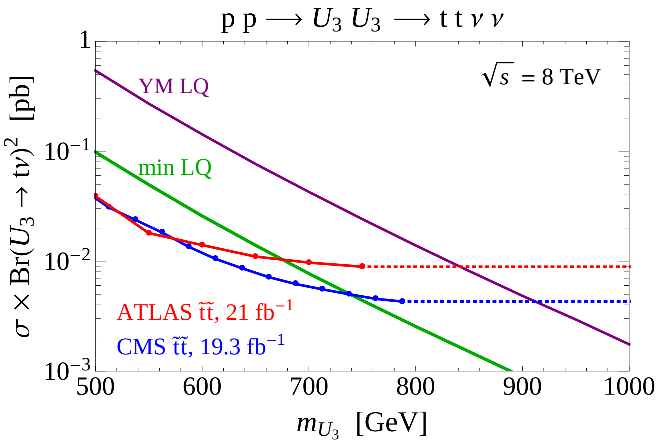

FIG. 4. Limits on m_{U_3} from direct $pp \to U_3 U_3$ production.

Direct leptoquark searches exist for the first decay mode, while bounds on the second mode can be derived by recasting stop searches. Vector leptoquark production at hadron colliders is complicated by the fact that in the absence of a UV completion, a marginal coupling not fixed by the minimal coupling prescription is present. In addition to the interactions given by Eq. (21), the rest of the Lagrangian for the vector leptoquark is

$$\mathcal{L}_U = -\frac{1}{2} U^\dagger_{\mu\nu} U^{\mu\nu} + m_U^2 U^\dagger_\mu U^\mu - i g_s \kappa\, (U^\dagger_\mu t^a U_\nu)\, G_a^{\mu\nu}, \quad (26)$$

up to higher-order operators suppressed by the cutoff. Here $U_{\mu\nu} = D_\mu U_\nu - D_\nu U_\mu$, while an additional gauge-invariant "dipole term" beyond the minimal coupling prescription is also present. If the vector leptoquark is a massive gauge boson from a spontaneously broken symmetry, then $\kappa = 1$, but in principle κ can be arbitrary. Setting $\kappa = 0$ gives the minimal production cross section. In all cases, a bound is produced directly on leptoquark masses, since the leptoquarks are dominantly produced due to their gauge couplings, and not their couplings to fermions. These are shown in Fig. 4. Since the leptoquark direct production rate is higher than that of stops, data up to the leptoquark exclusion limit is not publicly available (although stop searches could be recast for higher masses, which we leave for future work). In the figure we present a conservative extrapolation.

Additional bounds come from meson mixing and precision electroweak measurements of $Z \to \tau^+\tau^-$. These were computed in Ref. [77]. The interpretation of these is complicated by the fact that massive vector leptoquarks are not in themselves UV-complete and yield logarithmic divergences at loop level. In addition to $b \to s\nu\bar\nu$, the same interactions that lead to Eq. (25), also give rise to a quadratically divergent contribution to $B_{d,s}$ mixing at one-loop level. As a conservative estimate, the contribution of the finite part is then

$$\mathcal{O}_{\Delta B=2} = \frac{y_t^4 y_b^4 \lambda'^2 \lambda^2}{4\pi^2 m_{U_3}^2} |V_{tb}^* V_{ts}|^2\, (\bar b_L \gamma^\mu s_L)^2. \quad (27)$$

With current constraints [78], and again fixing $\lambda^2/m_{U_3}^2$ to explain the observed excess, this implies $\lambda'^2 \lesssim 0.1$, consistent with reasonable values of λ without any greater tuning than already necessary.

In the MFV framework, we thus find that two models are compatible with the data, once mild tuning is allowed for subleading MFV couplings to evade constraints from $b \to s\nu\bar\nu$: a scalar leptoquark in the $(\bar{\mathbf{3}},\mathbf{1})_{1/3}$ and a vector leptoquark in the $(\mathbf{3},\mathbf{3})_{2/3}$ representation of the SM gauge groups and in the (anti)fundamental of the $U(3)_d$ flavor symmetry group. Contrary to assumptions made elsewhere, although highly constrained, MFV physics can explain the $R(D^{(*)})$ data.

IV. SUMMARY AND FUTURE SIGNALS

We studied possible explanations of the observed enhancements of the $\bar B \to D^{(*)}\tau\bar\nu$ rates compared to the SM predictions. We identified which higher dimension operators provide good fits to the data and determined their required coefficients. Since a sizable modification of a tree-level process in the SM is required, the enhancements of $\bar B \to D^{(*)}\tau\bar\nu$ point to a light, tree-level, mediator. We found viable models, consistent with MFV, where the mediator is a leptoquark. While these models are consistent with present flavor and LHC data, the low mass scale of the mediator, $m \lesssim 1$ TeV, implies a variety of promising signals at future experiments.

With more precise future data from Belle II and LHCb, the necessary size of flavor violation may decrease while maintaining a high significance for NP. In this paper, we focused on NP that can accommodate the central values of the current measurements, but in reality NP may have a heavier mass scale, alleviating possible tensions with present constraints. If the magnitude of the deviation from the SM decreases while its experimental significance does not, it will become important to control theoretical uncertainties as well as possible. It would likely be advantageous to consider ratios [79] in which the range of q^2 integration is the same in the numerator and the denominator,

$$\widetilde{R}(D^{(*)}) = \frac{\displaystyle\int_{m_\tau^2}^{(m_B - m_{D^{(*)}})^2} \frac{d\Gamma(B \to D^{(*)}\tau\bar\nu)}{dq^2}\, dq^2}{\displaystyle\int_{m_\tau^2}^{(m_B - m_{D^{(*)}})^2} \frac{d\Gamma(B \to D^{(*)} l\bar\nu)}{dq^2}\, dq^2}, \quad (28)$$

which have smaller theoretical uncertainties than $R(D^{(*)})$. The reason is that including the $0 < q^2 < m_\tau^2$ region in the denominator simply dilutes this ratio and the sensitivity to new physics effects, and also because the uncertainties of the semileptonic form factors increase at smaller q^2 (larger recoil).

We conclude by enumerating future experimental measurements and new physics search channels that can test for models motivated by the measured $\bar B \to D^{(*)}\tau\bar\nu$ rates.

1. More precise data on $\bar{B} \to D^{(*)}\tau\nu$, including measurements of differential distributions and their correlations. These would also help discriminate NP from the SM, and between different NP models.
2. $D \to \pi\nu\bar{\nu}$ could be near the 10^{-5} level, which may be observable in BES III [80].
3. In some scenarios a nonresonant contribution to $t \to c\tau^+\tau^-$ is possible. While current $t \to cZ$ searches [81, 82] focus on $m_{\ell^+\ell^-}$ consistent with m_Z, there could also be a contribution that is inconsistent with the Z lineshape.
4. A NP contribution to $t \to b\tau\bar{\nu}$ of a few times $10^{-3}\,\Gamma_t$ is possible. Noting that the $\tau\bar{\nu}$ do not come from a W decay in the NP contribution, a high transverse mass tail, $m_T(\tau, \not{E}_T) > m_W$, can be searched for in semileptonic $t \to \tau$ decays.
5. In this paper we neglected CP violation. A simple experimental test is to constrain the difference between B^+ and B^- rates to $D^{(*)}\tau\bar{\nu}$; a measurable effect would be a clear sign of new physics.
6. Models that contribute to $\bar{B} \to D^{(*)}\tau\nu$ can modify $Z \to \tau^+\tau^-$ at the order of present constraints. Therefore, future precision Z-pole measurements may detect deviations from the SM.

Finally, we note that there may be viable models in which the $B \to D^{(*)}\tau\bar{\nu}$ rates are given by the SM, but $B \to D^{(*)}l\bar{\nu}$ ($l = e, \mu$) are suppressed by interference between NP and the SM [83].

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Appendix A: Comment on $U(2)^3$ scenarios

We presented several MFV models capable of describing the $R(D^{(*)})$ data while also coupling to other quark generations consistently with all constraints. As detailed earlier, extending the MFV framework to the lepton sector in these scenarios is not always possible, and so we have assumed horizontal symmetries governing the structure of lepton flavor where necessary. However, it is possible to treat quark and lepton flavor uniformly in all scenarios, at the cost of a slightly less predictive flavor framework.

A minimal self-consistent extension of the MFV framework exists in the form of the $U(2)^3$ flavor symmetry approach [84, 85]. Here, the UV physics is only assumed to preserve a $U(2)_Q \times U(2)_u \times U(2)_d$ symmetry for two generations in the quark sector, which is closely aligned with the first two generations of the SM. Three spurions are now necessary to minimally characterize the quark sector flavor symmetry breaking, here denoted as

$$\Delta_u \sim (\mathbf{2}, \bar{\mathbf{2}}, \mathbf{1})\,, \quad \Delta_d \sim (\mathbf{2}, \mathbf{1}, \bar{\mathbf{2}})\,, \quad \boldsymbol{V} \sim (\mathbf{2}, \mathbf{1}, \mathbf{1})\,, \quad \text{(A1)}$$

while the quark fields decompose as

$$\boldsymbol{q}_L \sim (2,1,1)\,, \quad \boldsymbol{u}_R \sim (1,2,1)\,, \quad \boldsymbol{d}_R \sim (1,1,2)\,, \quad \text{(A2)}$$

and the singlet third generation fields are q_{3L}, t_R, and b_R. The extension to the lepton sector, with an additional $U(2)_L^2 \times U(2)_e$ symmetry, is straightforward, and only a single spurion, $\Delta_e \sim (\mathbf{2}, \bar{\mathbf{2}})$ is needed (ignoring m_ν, which requires additional but much smaller spurions). Unlike the case of MFV, none of these spurions have $\mathcal{O}(1)$ components, so higher-order spurion insertions are highly suppressed.

Coupling predominantly to third generation quarks and leptons can now be imposed by demanding that the mediator is a singlet under all flavor symmetries, as long as currents of light quarks and leptons cannot be written as flavor singlets.

As an illustrative example, we consider a scalar leptoquark. Due to the V spurion, this leptoquark can couple to all quark generations. The interaction terms are

$$\delta\mathcal{L} = S\left(\lambda_1\, \bar{q}^c_{3L} i\tau_2 \ell_{3L} + \lambda_2 \boldsymbol{V}\, \bar{\boldsymbol{q}}^c i\tau_2 \ell_{3L} + \tilde{\lambda}_1\, \bar{t}^c_R \tau_R + \tilde{\lambda}_2 \boldsymbol{V}^\dagger \Delta_u \bar{\boldsymbol{u}}^c \tau_R\right). \quad \text{(A3)}$$

Coupling to the other lepton generations is impossible, since no combination of spurions can absorb their flavor charge. The contribution due to terms proportional to $\tilde{\lambda}_2$ can be neglected, as it always comes with multiple spurion suppression factors. Then the leading contribution by spurion counting to $b \to c$ transitions is given by

$$\frac{C''_{S_R}}{\Lambda^2} = \frac{V^*_{cb}}{m_S^2}\,\lambda_\alpha \lambda_1\,, \qquad \frac{C''_{S_L}}{\Lambda^2} = \frac{V^*_{cb}}{m_S^2}\,\lambda_\alpha \tilde{\lambda}_1\, y_c\,. \quad \text{(A4)}$$

Here λ_α is in the range (λ_1, λ_2) with the precise value set by the coefficients of $\boldsymbol{V}$ in up- and down-type Yukawa couplings. In the limit of $\lambda_1 = \lambda_2$ (and $\tilde{\lambda}_1 = \tilde{\lambda}_2$ if we include higher order terms in our Yukawa diagonalization), we recover the contribution of the third generation leptoquark of an MFV scenario. This is as we would expect. If we were to add a leptoquark in, say, the $(\mathbf{1}, \mathbf{1}, \mathbf{2})$ representation to the model above, setting all the suprion coefficients of terms with $\ell_{3L}/\boldsymbol{\ell}$ to λ and those with $\tau_R/\boldsymbol{e}_R$ to

$\tilde{\lambda}$, we precisely recover the particle content and spurion structure of the $(\mathbf{1},\mathbf{1},\bar{\mathbf{3}})$ MFV model.

The analysis of the constraints follows the MFV case. The treatment of contributions to $b \to s\nu\bar{\nu}$ is different, however. The term λ_α can be written, ignoring phases and using $|V_{cb}| \ll 1$,

$$\lambda_\alpha = \frac{\lambda_2 - x_t\lambda_1}{x_b - x_t}\,, \tag{A5}$$

where $x_{b,t}$ are free parameters defined in Ref. [84]. However, in this case contributions to $b \to s\nu\bar{\nu}$ are generated by the same interactions. Similar to Eq. (A5) above, the coefficient of this interaction can be written as

$$\frac{V_{cb}}{2}\,\lambda_1\frac{\lambda_2 - x_b\lambda_1}{x_b - x_t}, \tag{A6}$$

so one can reduce the rate of the interaction by tuning $x_b \approx \lambda_2/\lambda_1$ without affecting the $b \to c\tau\bar{\nu}$ rate. A slight tuning is then again necessary to avoid this constraint, although arising in a different set of couplings.

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