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With or without $U(2)$? Probing non-standard flavor and helicity structures in semileptonic $B$ decays

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A B S T R A C T
Motivated by the recent hints of lepton flavor universality violation observed in semileptonic $B$ decays, we analyze how to test flavor and helicity structures of the corresponding amplitudes in view of future data. We show that the general assumption that such non-standard effects are controlled by a $U(2)_q \times U(2)_c$ flavor symmetry, minimally broken as in the Standard Model Yukawa sector, leads to stringent predictions on leptonic and semileptonic $B$ decays. Future measurements of $\mathcal{R}_{\nu\ell}, \mathcal{R}_{\nu\tau}, \mathcal{B}(B_{c,s} \to \ell \nu), \mathcal{B}(\bar{B} \to \pi \ell \nu), \mathcal{B}(\bar{B} \to \pi \ell \ell), \mathcal{B}(B_{d,s} \to D^{(*)})$, as well as various polarization asymmetries in $\bar{B} \to D^{(*)} \tau \bar{\nu}$ decays, will allow to prove or falsify this general hypothesis independently of its dynamical origin.

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1. Introduction

Present data exhibit intriguing hints of violations of Lepton Flavor Universality (LFU) both in charged-current [1–5] and neutral-current [6–11] semileptonic $B$ decays. These hints can be well described employing Effective Field Theory (EFT) approaches (see [12–16] for the early attempts), whose main ingredients are the assumptions that New Physics (NP) affects predominantly semileptonic operators, and that it couples in a non-universal way to different fermion species. In particular, NP should have dominant couplings to third generation fermions and smaller, but non-negligible, couplings to second generation fermions. Interestingly enough, this non-trivial flavor structure resemble the hierarchies observed in the Standard Model (SM) Yukawa couplings, opening the possibility of a common explanation for the two phenomena.

An effective approach to address the question of a possible connection between the LFU anomalies and the SM Yukawa couplings and, more generally, to investigate the flavor structure of non-standard effects at low-energies, is the assumption of an appropriate flavor symmetry and a set of symmetry-breaking terms. The flavor symmetry does not need to be a fundamental property of the underlying theory, it could simply be an accidental low-energy property. Still, its use in the EFT context provides a very powerful organizing principle for a bottom-up reconstruction of the underlying dynamics.

In the context of the recent anomalies, the flavor symmetry that emerges as particularly suitable to describe the observed data is $U(2)_q \times U(2)_c$, which is a subset of the larger $U(2)^5$ proposed in [17–19] as a useful organizing principle to address the hierarchies in the SM Yukawa couplings and, at the same time, allow large NP effects in processes involving third-generation fermions (as expected in most attempts to address the electroweak hierarchy problem).

The scope of the present paper is a systematic investigation of the consequences of this symmetry hypothesis in (semi)leptonic $B$ decays. Contrary to previous analyses, where this symmetry has been implemented in the context of specific new-physics models, our goal here is to investigate the consequences of this symmetry (and symmetry-breaking) ansatz in general terms, with a minimal set of additional assumptions about the dynamical origin of the anomalies. As we will show, in the case of charged-current interactions, the symmetry ansatz alone is sufficient to derive an interesting series of testable predictions. The predictive power is smaller for neutral-current transitions, but also in that case we can identify a few clean predictions which are direct consequences of the symmetry ansatz alone.

2. The $U(2)^5$ symmetry in the SM

The $U(2)^5 \times U(1)_{B_3} \times U(1)_{L_3}$ symmetry is the global symmetry that the SM Lagrangian exhibits in the limit where we ne-
glect all entries in the Yukawa couplings but for third generation masses [17–19]. Under this symmetry, the first two SM fermion families transform as doublets of a given $U(2)$ subgroup,

$$U(2)^3 \equiv U(2)_l \times U(2)_r \times U(2)_u \times U(2)_d \times U(2)_e,$$  

while third-generation quarks (leptons) are only charged under $U(1)_{\beta_1}, \beta_3$. The largest breaking of this symmetry in the complete SM Lagrangian is controlled by the small parameter

$$\epsilon = \left[ \text{Tr}(Y_u Y_d^\dagger) - \text{Tr}(Y_u' Y_d'^\dagger) \right] / \text{Tr}(Y_d Y_d^\dagger) \approx 0.04. \tag{2}$$

A minimal set of $U(2)^5$ breaking terms (spurions) which let us reproduce all the observable SM flavor parameters (in the limit of vanishing masses), without tuning and with minimal size for the breaking terms, is

$$V_q \sim (2, 1, 1, 1, 1), \quad V_\ell \sim (1, 2, 1, 1, 1), \quad \Delta_u(d) \sim (2, 1, 2, 1, 1, 1, 1, 1). \tag{3}$$

In terms of these spurions, the $3 \times 3$ Yukawa matrices can be decomposed as

$$Y_{u(d)} = Y_{\text{tr}}(\Delta_{u(d)} \times |\epsilon| Y_q), \quad Y_\ell = Y_{\text{tr}}(\Delta_{u(d)} \times |\epsilon| Y_\ell), \tag{4}$$

where $x_{\text{tr}, \ell}$ and $y_{\text{tr}, \ell}$ are free complex parameters, expected to be of $O(1)$, Note that $\Delta_{u(d)}$ are $2 \times 2$ complex matrices, while $Y_{u(d)}$ are 2-dimensional complex vectors.

The precise size of the spurions is not known; however, we can estimate it by the requirement of no tuning in the $O(1)$ parameters. This implies $|V_q| = O(\epsilon)$. In the limit of vanishing neutrino masses, the size of $|V_\ell|$ cannot be unambiguously determined. As discussed below (see also [20]), a good fit of the anomalies in semileptonic $B$ decays is obtained for

$$|V_L|, \quad |V_q| = O(10^{-1}), \tag{5}$$

which is perfectly consistent with: i) the estimate $|V_q| = O(\epsilon)$; ii) the hypothesis of a common origin for the two leading $U(2)^5$ breaking terms in quark and lepton sectors. The entries in the $2 \times 2$ matrices $\Delta_{u(d)}$ are significantly smaller than $|V_q|$, with a maximal size of $O(10^{-2})$ in the quark sector.

By appropriate field redefinitions and without loss of generality, one can remove unphysical parameters in the Yukawa matrices in (4) (see App. A). Working in the so-called interaction basis, where the second generation in $U(2)_q$ space is defined by the alignment of the leading spurious,

$$V_q = |V_q| \times \tilde{n}, \quad \tilde{n} = \begin{pmatrix} 0 \\ 1 \end{pmatrix}, \tag{6}$$

one can bring the Yukawa matrices to the following form

$$Y_u = |y_r| \begin{pmatrix} Y_u' \tilde{\Delta}_u & |V_q| e^{i\phi_q} \tilde{n} \\ 0 & 1 \end{pmatrix}, \tag{7}$$

$$Y_d = |y_b| \begin{pmatrix} Y_d' \tilde{\Delta}_d & |V_q| e^{i\phi_q} \tilde{n} \\ 0 & 1 \end{pmatrix}, \tag{7}$$

where all $\tilde{\Delta}_{u(d)}$ are $2 \times 2$ diagonal matrices, $O_{u(e)}$ are $2 \times 2$ orthogonal matrices and $U_q$ is of the form

$$U_q = \begin{pmatrix} c_{\phi} & -s_{\phi} \exp(i\theta_q) \\ s_{\phi} & c_{\phi} \end{pmatrix}, \tag{8}$$

with $s_{\phi} = \sin \theta_{\phi}$ and $c_{\phi} = \cos \theta_{\phi}$.

The Yukawa matrices in (7) get diagonalized by means of appropriate unitary transformations: $U_q^\dagger Y_q R_f = \text{diag}(Y_q)$, with $f = u, d, e$. The most general form for these unitary transformations is

$$L_d \approx \begin{pmatrix} c_{\phi} & -s_{\phi} \exp(i\theta_q) & 0 \\ s_{\phi} & c_{\phi} \exp(i\theta_q) & 0 \\ 0 & 0 & s_{\phi} \end{pmatrix}, \tag{8}$$

$$R_d \approx \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & -m_{\tilde{s}} & s_{\tilde{s}} \end{pmatrix}, \tag{9}$$

$$L_e \approx \begin{pmatrix} c_{\tau} & -s_{\tau} & 0 \\ s_{\tau} & c_{\tau} & 0 \\ 0 & 0 & c_{\tilde{s}} \end{pmatrix}, \tag{9}$$

$$R_e \approx \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & -s_{\tilde{s}} & c_{\tilde{s}} \end{pmatrix}, \tag{9}$$

with $L_q = U_q V_{\text{CKM}}$. Here we have taken advantage of the constraints imposed by fermions masses and CKM matrix elements to eliminate various parameters appearing in $L_f$ and $R_f$.

The further imply that $s_{\phi}$ and $c_{\phi}$ are constrained by $s_{\phi}/c_{\phi} = |V_{u(d)}/V_{13}|$ and $c_{\phi} = \text{arg}(V_{12}^\ast V_{e3})$, and that $s_{\phi} = s_{\phi} - V_{e3}$. The light-family leptonic mixing ($s_{\phi}$), appearing in $O_{\tau\ell}$ cannot be expressed in terms of measurable quantities. Two additional mixing angles which remain un constrained are $s_{\tilde{g}}/c_{\tilde{g}} = |V_{u(d)}/V_{13}|$ and $s_{\tilde{g}}/c_{\tilde{g}} = |V_{u(d)}/V_{13}|$. Finally, $\phi_q$ is an unconstrained $O(1)$ phase, that becomes unphysical in the limit $s_{\phi} \to 0$ (or, equivalently, $s_{\tilde{g}} \to 0$). This limit is phenomenologically required in models where $\Delta F = 2$ operators are generated at tree-level around the TeV scale: in such case one needs to impose a mild alignment to the down basis, i.e. $|s_{\phi}| \lesssim 0.1 e$, to satisfy the constraints from $B_s, d$-meson mixing [21–23].

3. Impact of $U(2)^5$ on the EFT for semileptonic $B$ decays

Having defined the flavor symmetry and its symmetry breaking terms from the SM Yukawa sector, we are ready to analyze its implications beyond the SM. Assuming no new degrees of freedom below the electroweak scale, we can describe NP effects in full generality employing the so-called SMEFT. We limit the attention to dimension-six four-fermion operators bilinear in quark and lepton fields,\(^2\) that we write generically as

$$\mathcal{L}_{\text{EFT}} = -\frac{1}{2 \eta} \sum_{k,\beta_1,\beta_2} c_k^{(\beta_1,\beta_2)} c_k^{(\beta_1,\beta_2)} + \text{h.c.}, \tag{10}$$

\(^2\) The removal of unphysical parameters presented in App. A corrects a similar analysis presented in [21], where it was erroneously concluded that the parameter $\alpha_{\phi}$ is unconstrained.

\(^3\) We neglect operators which modify the effective couplings of $W$ and $Z$ bosons. These are highly constrained and cannot induce sizable LFV violating effects.
where $v \approx 246$ GeV is the SM Higgs vev, $\{\alpha, \beta\}$ are lepton-flavor indices, and $\{i, j\}$ are quark-flavor indices. The operators in the Warsaw basis [24] with a non-vanishing tree-level matrix element in semileptonic $B$ decays are

\[
\begin{align*}
O^{(1)}_{\ell q} &= (\bar{q}_L^i \gamma^\mu \ell_L^j) (\bar{q}_L^i \gamma_\nu q_L^j), \\
O^{(3)}_{\ell q} &= (\bar{q}_L^i \gamma^\mu \tau^a \ell_L^j) (\bar{q}_L^i \gamma_\nu \ell_L^j), \\
O_{\ell d} &= (\bar{q}_L^i \gamma^\mu \ell_L^j) (\bar{q}_R^i \gamma_\nu d_R^j), \\
O_{\ell e} &= (\bar{q}_L^i \gamma^\mu \ell_L^j) (\bar{q}_L^i \gamma_\nu \ell_L^j), \\
O_{\ell d \ell d} &= (\bar{q}_L^i \gamma^\mu \ell_L^j) (\bar{q}_R^i \gamma_\nu d_R^j), \\
O_{\ell d \ell e} &= (\bar{q}_L^i \gamma^\mu \ell_L^j) (\bar{q}_L^i \gamma_\nu \ell_L^j), \\
O^{(1)}_{\ell e \ell q} &= (\bar{q}_L^i \gamma^\mu \ell_L^j) (\bar{q}_L^i \gamma_\nu \ell_L^j), \\
O^{(3)}_{\ell e \ell q} &= (\bar{q}_L^i \gamma^\mu \ell_L^j) (\bar{q}_L^i \gamma_\nu \ell_L^j), \\
O_{\ell d \ell e} &= (\bar{q}_L^i \gamma^\mu \ell_L^j) (\bar{q}_R^i \gamma_\nu \ell_L^j),
\end{align*}
\]  

(11)

where $\tau^i$ are the Pauli matrices and $\{a, b\}$ are $SU(2)_L$ indices. Our main hypothesis is to reduce the number of $C^{ij}_{1,2}$ retaining only those corresponding to $U(2)^3$ invariant operators, up to the insertion of one or two powers of the leading $U(2)_q \times U(2)_e$ spurions in (5).

A first strong simplification arises by neglecting subleading spurions with non-trivial transformation properties under $U(2)_{d,e}$. Since we are interested in processes of the type $b \to c(u)\ell$ and $b \to s(d)\ell$, this implies that only the operators $O^{(1)}_{\ell q}$, $O^{(3)}_{\ell q}$, $O_{\ell d e}$ and $O_{\ell d \ell d}$ can yield a relevant contribution. Among those, $O_{\ell d e}$ can significantly contribute at tree-level only to $b \to s\tau\ell$ transitions; since the latter are currently poorly constrained [see sect. 4.3], we do not consider this operator for simplicity. We are thus left with the following effective Lagrangian

\[
\mathcal{L}_{\text{EFT}} = -\frac{1}{v^2} \left[ C_{V_1} \Lambda^{[ij]ab}_{V_1} O^{(1)}_{\ell q} + C_{V_3} \Lambda^{[ij]ab}_{V_3} O^{(3)}_{\ell q} + 2 C_S \Lambda^{[ab]ij}_{S} (O_{\ell d e} + \text{h.c.}) \right],
\]

(12)

where $C_{V_1,3}$ control the overall strength of the NP effects and $\Lambda_{V, S}$ are tensors that parametrize the flavor structure. They are normalized by setting $\Lambda^{V}_{V, S} = 1$, which is the only term surviving in the exact $U(2)^3$ limit.

Let us consider first the structure of $\Lambda^{[ij]ab}_{S}$, which is particularly simple. Neglecting $U(2)_{d,e}$ breaking spurions, it factorizes to

\[
\Lambda^{[ij]ab}_{S} = (\Gamma^j_L)^{ij} \times \Gamma^R_L,
\]

(13)

where, in the interaction basis,

\[
\Gamma^a_L = \begin{pmatrix} x_{\ell q} (V^j_L)^* & x_{\ell q} V^j_L \\ x_{\ell q} (V^j_L)^* & 1 \end{pmatrix}, \quad \Gamma^R_L = \begin{pmatrix} 0 & 0 \\ 0 & 1 \end{pmatrix}.
\]

(14)

Here $x_{\ell q,L}$ are $O(1)$ coefficients and we have neglected higher-order terms in $V_{\ell q}$ (that would simply redefine such coefficients). Moving to the mass-eigenstate basis of down quarks and charged leptons, where

\[
q_L^i = \begin{pmatrix} V_L^{\nu} u_L^i \\ d_L^i \end{pmatrix}, \quad \nu^a = \begin{pmatrix} 1 \\ 0 \end{pmatrix},
\]

(15)

we have $\Gamma_L \to \hat{\Gamma}_L \equiv L^T \Gamma L$ and $\Gamma_R \to \hat{\Gamma}_R \equiv R^T \Gamma_R R_e$ [see (9)], with the new matrices assuming the following explicit form in $3 \times 3$ notation

\[
\hat{\Gamma}_L = e^{i \phi_0} \begin{pmatrix} \Delta^{\mu\nu}_{q_L} & 0 & \Delta^{\mu\nu}_{q_R} \\ \Delta^{\mu\nu}_{q_L} & \lambda^\mu_{L} & \lambda^\nu_{L} \\ \Delta^{\mu\nu}_{q_R} & \lambda^\nu_{R} & \lambda^\mu_{R} \end{pmatrix}, \quad \hat{\Gamma}_R \approx e^{i \phi_0} \begin{pmatrix} 0 & 0 & \lambda^\mu_{L} \\ 0 & 0 & \lambda^\nu_{L} \\ -m_e/m_\tau & 1 \end{pmatrix},
\]

(16)

The (complex) parameters $\lambda^\mu_{L, R}$, $\lambda^\nu_{L, R}$, and $\Delta^{\mu\nu}_{q_R}$ are a combination of the spurions in (14) and the rotation terms from $L_{d,e}$, that satisfy

\[
\lambda^\mu_{L} = O(|V_{qL}|), \quad \lambda^\mu_{R} = O(|V_{qR}|), \quad \lambda^\nu_{L} = O(1), \quad \Delta_{qR}^{\mu\nu} = O(\lambda^\mu_{L, R}),
\]

\[
\frac{\lambda^\mu_{L}}{\lambda^\nu_{qL}} = \frac{\Delta_{qR}^{\mu\nu}}{\lambda^\nu_{qL}} = V_{qL}^\nu, \quad \frac{\lambda^\nu_{L}}{\lambda^\mu_{qL}} = \frac{\Delta_{qR}^{\mu\nu}}{\lambda^\mu_{qL}} = s.e.
\]

(17)

On the r.h.s. of the first line of (16) we have neglected tiny terms suppressed by more than two powers of $|V_{qL, R}|$ or $s.e.$

If we consider at most one power of $V_{q_L}$ and one power of $V_{q_R}$, then also $\Lambda^{[ij]ab}_{V_1}$ factorizes into

\[
\Lambda^{[ij]ab}_{V_1} = (\Gamma^L_{V_1})^a \gamma^j \times (\Gamma^L_{V_1})^b,
\]

(18)

where $\Gamma^L_{V_1}$ and $\Gamma^L_{V_3}$ have the same structure as $\Gamma_L$ with a priori different $O(1)$ coefficients for the spurions. Moving to the basis (15), $\Gamma_{V_1}^j$ assumes the same structure as $\hat{\Gamma}_L$ in (16), with parameters which can differ by O(1) overall factors, but that obey the same flavor ratios as in (17). Corrections to the factorized structure in (18) arises only to second order in $V_{q_L}$ or $V_{q_R}$, generating terms which are either irrelevant or can be reabsorbed in a redefinition of the observable parameters in the processes we are interested in (see sect. 4).

3.1. Matching to the $U_1$ leptoquark case

The EFT in (12), with factorized flavor couplings as in (13) and (18), nicely matches the structure generated by integrating out a $U_1$ vector leptoquark, transforming as $(3, 1/2)_3$ under the SM gauge group. As noted first in [16], this field provides indeed an excellent mediator to build in a natural, and sufficiently general way, an EFT for semileptonic $B$ decays built on the $U(2)^3$ flavor symmetry broken only by the leading $V_{q_L}$ and $V_{q_R}$ spurions (see [20–23, 25–30] for other phenomenological analysis of the $U_1$ leptoquark in B physics).

Writing the interaction between the $U_1$ field and SM fermions in the basis of (15) as

\[
\mathcal{L}_{U_1} = \frac{SU}{\sqrt{2}} \left[ \beta_{1L} \bar{q}_L^i (\bar{q}_L^i \gamma_\mu \ell_L^j) + \beta_{1R} (\bar{q}_R^i \gamma_\mu \ell_L^j) \right] U_1^\mu + \text{h.c.},
\]

(19)

the flavor symmetry hypothesis imply a parametric structure for $\beta_{1L}$ and $\beta_{1R}$ identical to that of $\Gamma^L_{V_1}$ and $\Gamma^L_{V_3}$ in (16). Normalizing $g_{U_1}$ such that $\beta_{1L} = 1$, and integrating out the $U_1^\mu$ field, leads to the following (tree-level) matching conditions for the parameters of $\mathcal{L}_{\text{EFT}}$

\[
\Gamma_{V_1}^{ij} = \Gamma_{L}^{ij} = \Gamma_{R}^{ij}, \quad C_V = C_{V_1} = C_{V_3} = \frac{g_{U_1}^2 v^2}{4 M_U} > 0,
\]

(20)

and

\[
\frac{C_S}{C_V} = -2 \beta_R \frac{\lambda^\mu_{L}}{\lambda^\nu_{qL}} - \beta_{1L}^\mu - \beta_{1L}^\nu - \lambda_{qR}^{\mu\nu} = -2 \beta_R \frac{\lambda^\mu_{L}}{\lambda^\nu_{qL}}
\]

(21)
where \( \beta_R \equiv \beta_R^{\nu \tau} \). Note that while \((21)\) is just a redefinition of the free parameters of the effective Lagrangian, \((20)\) give non-trivial constraints. We also stress that, beside the overall coupling (encoded in \(C_V\)) the four combinations of couplings in \((21)\) indicate the helicity structure of the interactions \((C_S/C_V)\) and its alignment in quark and lepton flavor space \((\lambda^q, \lambda^l_i, \text{and } \Delta^q l)\).

The condition \(C_V = C_{V_1}\), arising naturally in the \(U_1\) case, is important to evade the tight constraints on \(C^V_Q\) from \(b \to s V\bar{v}\) observables and electroweak precision tests \([27]\). In order to analyze charged currents as much independent as possible from other observables, we take \(C_V \equiv C_{V_1} = C_{V_1}\) in the Lagrangian in \((12)\). Similarly, we set \(\Gamma_{\tau_1}^{V} = \Gamma_{\tau_1}^{V} = \Gamma_{\tau_1}^L\) in order to avoid the introduction of redundant parameters as far as neutral currents are concerned.

4. The relevant observables

4.1. \(b \to c \tau \bar{v}\) rates and polarization asymmetries

At fixed quark and lepton flavors, the effective Lagrangian in \((12)\) depends only on two coefficients. In the \(b \to c \tau \bar{v}\) case, we conveniently re-define them as

\[
C^V_{V(S)} = C^V_{V(S)} \left[ 1 + \frac{\lambda_\tau^q}{\lambda_{\tau_1}^q} \left( \frac{V_{cs}}{V_{cb}} + \frac{V_{cd}^*}{V_{td}} \right) \right] = C^V_{V(S)} \left[ 1 - \frac{\lambda_\tau^q}{\lambda_{\tau_1}^q} \left( \frac{V_{ts}^*}{V_{tb}} \right) \right],
\]

(22)

where, in the last line, we have used CKM unitarity. When defining \(C^V_{V(D)}\), we have factorized the CKM factor \(V_{cb}\), such that the left-handed part of the interactions is modified as

\[
A_{SM}^{\nu \tau} \to (1 + C^V_{V}) A_{SM}^{\nu \tau}.
\]

(23)

In the absence of the simplifying hypothesis \(\Gamma_{\tau_1}^{V_1} = \Gamma_{\tau_1}^L\), one would need to redefine \(C^V_{V}\) replacing \(\lambda_\tau^q\) with \(\lambda_{\tau_1}^q\). Employing this hypothesis, as in the leptoquark case, the ratio \(C^V_{S}/C^V_{V} = C_S/C_V\) is flavor blind and depends only on the helicity structure of the NP amplitude.

Using the results in \([32,33]\) for the \(\bar{b} \to D^{(*)} l \bar{v}\) form factors and decay rates, and neglecting the tiny NP corrections in the \(\ell = \mu, e\) case (see below) leads to the following expression for the LRU ratios, \(R_L = \Gamma(\bar{b} \to H \tau \bar{v})/\Gamma(\bar{b} \to H l \bar{v})\),

\[
\frac{R_{D}}{R_{D}^{SM}} \approx |1 + C^V_{S}|^2 + 1.50(1) \text{Re}[(1 + C^V_{V}) \eta_S C^*_S] + 1.03(1) |\eta_S C^*_S|^2,
\]

\[
\frac{R_{D}^{*}}{R_{D}^{*}^{SM}} \approx |1 + C^V_{S}|^2 + 0.12(1) \text{Re}[(1 + C^V_{V}) \eta_S C^*_S] + 0.04(1) |\eta_S C^*_S|^2,
\]

(24)

where \(\eta_S \approx 1.7\) arises from the running of the scalar operator from the TeV scale down to \(m_b\) \([34]\). Updated SM predictions for \(R_{D}^{(*)}\) can be found in \([35]\):

\[
R_{D}^{SM} = 0.297(3), \quad R_{D}^{SM} = 0.250(3).
\]

(25)

Current measurements of \(R_D\) and \(R_{D}^{*}\) \([1-5]\) lead to the constraints on \(C^V_{S}\) and \(C^V_{V}\) shown in Fig. 1 (dashed contour lines) where, for simplicity, we have assumed these couplings to be real. For comparison, the directions corresponding to a pure left-handed \((\beta_R = 0)\) or a vector-like interaction \((\beta_R = -1)\) for the \(U_1\) are also indicated.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig1.png}
\caption{Best fit regions in the \((C^V_{S}, C^V_{V})\) plane. The three contours corresponds to 1, 2, and 3\(\sigma\) intervals. The dashed blue lines take into account only the information from \(R_D\) and \(R_{D}^{*}\) (for which we use the HFLAV average \([31]\), whereas the continuous lines also include constraints from \(b \to u\) observables (see sect. 4.2 for more details). The colored bands correspond to the 1\(\sigma\) regions defined by each observable. The gray bands show the 90\% CL exclusion region from \(B(\bar{b} \to \tau l f)\) (see sect. 4.3).}
\end{figure}

Before discussing the impact of additional observables in constraining the same set of parameters (under additional assumptions), we stress that once \(C^V_{S}\) and \(C^V_{V}\) have been determined, all the other \(b \to c \tau \bar{v}\) observables are completely fixed by the \(U(2)\) invariant structure of \(L_E\) and can be used to test it. Particularly interesting in this respect are the polarization asymmetries,

\[
F_{L}^{D} = \frac{\Gamma(\bar{b} \to D^{(*)} l \bar{v})}{\Gamma(\bar{b} \to D^{(*)} \tau \bar{v})},
\]

(26)

\[
F_{L}^{D^{*}} = \frac{1}{\Gamma(\bar{b} \to D^{(*)} l \bar{v})} \left( \frac{1}{\Gamma(\bar{b} \to D^{(*)} \tau \bar{v})} \right) = \frac{\Gamma(\bar{b} \to D^{(*)} l \bar{v})}{\Gamma(\bar{b} \to D^{(*)} \tau \bar{v})} + \frac{1}{\Gamma(\bar{b} \to D^{(*)} \tau \bar{v})};
\]

where the \(l^{(*)}\) denotes a \(l\) with \(\pm 1/2\) helicity. We find the following expressions for these observables\(^4\)

\[
\frac{F_{L}^{D^{*}}}{F_{L}^{D^{*} SM}} \approx \left( \frac{R_{D}^{*}}{R_{D}^{SM}} \right)^{-1} \left( |1 + C^V_{S}|^2 + 0.087(4) |\eta_S C^*_S|^2 + 0.253(8) \text{Re}[(1 + C^V_{V}) \eta_S C^*_S] \right),
\]

\[
\frac{p_{F}^{D}}{p_{F}^{D^{*} SM}} \approx \left( \frac{R_{D}^{*}}{R_{D}^{SM}} \right)^{-1} \left( |1 + C^V_{S}|^2 + 3.24(1) |\eta_S C^*_S|^2 + 4.69(2) \text{Re}[(1 + C^V_{V}) \eta_S C^*_S] \right),
\]

(27)

\[
\frac{p_{F}^{D^{*}}}{p_{F}^{D^{*} SM}} \approx \left( \frac{R_{D}^{*}}{R_{D}^{SM}} \right)^{-1} \left( |1 + C^V_{S}|^2 - 0.079(5) |\eta_S C^*_S|^2 - 0.23(1) \text{Re}[(1 + C^V_{V}) \eta_S C^*_S] \right).
\]

\(^4\) The numerical coefficients in \([27]\) are compatible with those obtained in \([36]\), within the errors, and are within 5\% of to those obtained in \([37]\) (where different form factors have been employed).
where $B(\bar{B}_c \to \tau \bar{\nu}_\tau)_{SM} \approx 0.02$, $\chi_c = m_{\bar{B}_c}/(m_{\tau} + m_{\bar{\nu}_\tau}) \approx 4.3$. We find $B(\bar{B}_c \to \tau \nu)$ to be at most at the level of 10% for the best fit contours in Fig. 1, well below the $B(\bar{B}_c \to \tau \bar{\nu}) \leq 30\%$ bound obtained in [41].

In principle, additional probes of the $b \to c \tau \bar{\nu}$ amplitude are provided by the $\tau/\mu$ LFU ratios in $\Lambda_{b} \to \Lambda_c \ell \bar{\nu}_\ell$ [37] and in $\bar{B}_c \to \psi \ell \bar{\nu}$ [43] decays. In both cases scalar amplitudes are subleading and, in our framework, one should expect an enhancement compared to the SM prediction similar to the one occurring in $R_{D^*}$. However, measuring the LFU ratio in $\Lambda_{b} \to \Lambda_c \ell \bar{\nu}$—where we have a precise SM prediction [44]—is quite challenging, while in the $\bar{B}_c \to \psi \ell \bar{\nu}$ case the current SM theory error is well above the 10% level [45].

4.2. $b \to u \tau \bar{\nu}$ transitions

The analog of $C_{\nu}(S)$ for $b \to u$ transitions are the effective couplings

$$C_{\nu}(S) = C_{\nu}(S) \left[ 1 + \lambda^2 \left( \frac{V_{ub}}{V_{ub}} + \frac{V_{ub}}{V_{ub}} \frac{V_{ud}}{V_{td}} \right) \right] = C_{\nu}(S),$$

(32)

where the result in the second line follows from CKM unitarity. The prediction of same size NP effects, relative to the SM, in $b \to u$ and $b \to c$ transitions is a distinctive feature of the minimally-broken $U(2)^3$ hypothesis.

At present, the most significant constraints on $b \to u \tau \bar{\nu}$ transitions are derived from $\bar{B}_u \to \tau \nu$, whose branching ratio is

$$\frac{B(\bar{B}_u \to \tau \nu)}{B(\bar{B}_u \to \tau \nu)_{SM}} = \left[ 1 + C_{\nu} + \chi_u \eta_{S} C_{S}^2 \right],$$

(33)

where $\chi_u = m_{\bar{B}_u}^2/(m_{\tau} + m_{\bar{\nu}_\tau}) \approx 3.75$ and $B(\bar{B}_u \to \tau \nu)_{SM} = 0.812(54)$ [46]. The continuous contour lines in Fig. 1 show the constraints on $C_{\nu}^2 = C_1^2$ and $C_S^2$ once the experimental data on $R_{D^*}$, are combined with those on $B(\bar{B}_u \to \tau \nu)$ [47] (1σ range indicated by the purple band in Fig. 1).

In the future, very interesting constraints are expected from $\bar{B} \to \pi \nu \bar{\nu}$. Using the hadronic parameters in [48,49], we find for $R_{\pi} \equiv B(\bar{B} \to \pi \nu \bar{\nu})/B(\bar{B} \to \pi \ell \bar{\nu})$

$$\frac{R_{\pi}}{R_{\pi}^{SM}} = \left[ 1 + C_{\nu}^2 + 1.13(7) \Re \left(1 + C_{\nu}^2 \right) \eta_{S} C_{S}^2 \right] + 1.36(9) \eta_{S} C_{S}^2,$$

(34)

where $R_{\pi}^{SM} = 0.641(16)$ [48,50,51]. In the limit where quadratic NP effects can be neglected, the following approximate relation holds

$$\frac{R_{\pi}^{SM}}{R_{\pi}^{SM}} \approx 0.75 \frac{R_{D}^{SM}}{R_{D}^{SM}} + 0.25 \frac{R_{D^*}^{SM}}{R_{D^*}^{SM}},$$

(35)

which would allow a non-trivial test of the $U(2)^3$ structure of the interactions. In Fig. 3 we show the predictions for $B(\bar{B}_u \to \tau \nu)$, $B(\bar{B} \to \pi \nu \bar{\nu})$, and $B(\bar{B}_c \to \tau \nu)$, as a function of $\Delta R_D - \Delta R_{D^*}$. As shown in the figure, our setup predicts also

$$\frac{B(\bar{B}_u \to \tau \nu)}{B(\bar{B}_u \to \tau \nu)_{SM}} \approx \frac{B(\bar{B}_c \to \tau \nu)}{B(\bar{B}_c \to \tau \nu)_{SM}},$$

(36)

where the difference among the two modes arises by sub-leading spectator mass effects in the chirality-enhancement factors $\chi_c$ and $\chi_u$. 

---

Fig. 2. Deviations of the polarization asymmetries compared to the SM as a function of $\Delta R_D - \Delta R_{D^*}$. The predictions are obtained using the fit in Fig. 1 (continuous lines). In gray, the experimental value of $\Delta F_D^{SM}$ at 1σ and 2σ.
4.3. $b \to s\ell\bar{\ell}(\nu\bar{\nu})$ transitions

The $b \to s$ semileptonic transitions have a rich phenomenology and have been extensively discussed in the recent literature. Contrary to the charged-current case, here model-dependent assumptions, such as the constraints in (20), play a more important role. Rather than presenting a comprehensive analysis of the various observables accessible in these modes, our scope here is to focus on: i) model-independent predictions related to the minimally-broken $U(2)^3$ hypothesis; ii) clean observables controlling the size of the symmetry breaking terms.

$b \to s\tau\bar{\tau}(\nu\bar{\nu})$. Under the assumption $C_{V_1} = C_{V_2}$, following from the hypothesis of a $U_1$ UV-completion, NP effects in $b \to s\nu\bar{\nu}$ transitions are forbidden at tree level. On the other hand, NP contributions to $b \to s\tau\bar{\tau}$ are almost as large as those in $b \to c\tau\bar{\nu}$ for $\lambda_s^a = 0.01$ (see e.g. [27,52]). The most relevant observable involving these transitions is $B(B_s \to \tau\bar{\tau})$, which could receive a sizable chiral enhancement:

$$\frac{B(B_s \to \tau\bar{\tau})}{B(B_s \to \tau\bar{\nu})_{SM}} = 1 + \frac{2\pi \lambda_q^s}{\alpha V_{tb} V_{ts}C^{10}_{10}} (C_9 + \chi_s \eta_5 C_S)^2$$

$$+ \left( 1 - \frac{4m_{B_s}^2}{m_{B_s}^2} \right)^2 \frac{2\pi \lambda_q^s}{\alpha V_{tb} V_{ts}C^{10}_{10}} \chi_S \eta_5 C_S^2,$$  

(37)

where $\chi_s = m_{B_s}^2/[2m_t(m_b + m_t)] \approx 1.9$ and $C^{10}_{10} \approx -4.3$ (see below). The enhancement of this rate compared to the SM expectation could reach a factor of $10^2$ ($10^3$) for $C_S = 0$ ($C_S = 2C_V$). However, the current experimental limit [53,54]

$$\frac{B(B_s \to \tau\bar{\tau})}{B(B_s \to \tau\bar{\nu})_{SM}} < 8.8 \times 10^3$$  

(95% CL),

is still well below the possible maximal enhancement. As a result, at present this observable does not put stringent constraints on the parameter space of the EFT: in Fig. 1 we show the 90% CL exclusion region in the $(C_S^2, C_V^2)$ plane for $\lambda_s^a = 3 |V_{ts}|$.

As pointed out in [55], a possible large enhancement of the $b \to s\tau\bar{\tau}$ amplitude can indirectly be tested via the one-loop-induced lepton-universal contributions to $b \to s\ell\bar{\ell}$ ($\ell = e, \mu, \tau$) in the $O_3$ direction (see below). This contribution is well compatible and even favored by current data [30,56].

$b \to s\mu\bar{\mu}(e\bar{e})$ FCNC decays to light leptons offer an excellent probe of the $U(2)^3$ breaking terms in the lepton sector. These transitions are commonly described in terms of the so-called weak effective Hamiltonian [57,58]

$$\mathcal{H}_{\text{weak}}^{1,2} \supset -\frac{4G_F}{\sqrt{2}} \frac{\alpha}{4\pi} V_{tb}V_{ts}^* \sum_{i=9,10,5,\mu} C_i^\mu O_i^\mu,$$  

(39)

with $G_F$ the Fermi constant, $\alpha$ the fine-structure constant and

$$C_9^\mu = (\tilde{G}_{\mu b} P_L b)(\tilde{\bar{\nu}}_\mu \tilde{\ell}), \quad O_9^\mu = (\tilde{G}_{\mu b} P_L b)(\tilde{\bar{\nu}}_\mu \tilde{\ell}),$$

$$C_5^\mu = (\tilde{S}_{\mu b} b)(\tilde{\bar{\nu}}_\mu \tilde{\ell}), \quad O_5^\mu = (\tilde{S}_{\mu b} b)(\tilde{\bar{\nu}}_\mu \tilde{\ell}).$$

(40)

In the SM, $C_9^\mu = 4.1$, $C_9^\mu \approx -4.3$ and $C_5^\mu = C_5^\mu = 0$. Matching to the Lagrangian in (12), we get ($C_i = C_i^{\text{SM}} + \Delta C_i$)

$$\Delta C_9^\mu = -\Delta C_9^{\mu*} = -\frac{2\pi}{\alpha V_{tb} V_{ts}^*} C_V^{\mu*} \frac{\lambda_q^\mu}{\lambda_q^*} \frac{\lambda_q^s}{\lambda_q^*} \frac{\lambda_q^0}{\lambda_q^*},$$  

$$C_9^\mu = -C_9^{\mu*} = -\frac{2\pi}{\alpha V_{tb} V_{ts}^*} C_V^{\mu*} \frac{\lambda_q^\mu}{\lambda_q^*} \frac{\lambda_q^s}{\lambda_q^*} \frac{\lambda_q^0}{\lambda_q^*},$$

(41)

while the corresponding tree-level effects in the electron sector are negligible.

One of most relevant observables involving these transitions are the LFU ratios $R_{K^{(*)}} = \Gamma(B \to K^{(*)}\mu\bar{\mu})/\Gamma(B \to K^{(*)}e\bar{e})$, which are particular interesting due to their robust theoretical predictions: $R_{K^{(*)}}^{\text{SM}} = 1.00 \pm 0.01$ [59]. In our setup, one gets [60,61]

$$R_K \approx R_{K^*} \approx 1 + 0.47 \Delta C_9^{\mu*}. $$

(42)

The prediction $R_K \approx R_{K^*}$, is a direct consequence of our flavor symmetry assumptions and is independent of the initial set of dimension-six SMEFT operators. As observed first in [62], the relation $R_K \approx R_{K^*}$ holds in any NP model where LFU contributions to $b\to s\ell\bar{\ell}$ decays are induced by a left-handed quark current: in our framework this is a direct consequence of the smallness of the flavor-symmetry breaking terms in the right-handed sector. From the experimental point of view, this implies that all $\mu/e$ universality ratios in $b \to s$ transitions are expected to be the same, provided their SM contribution is dominated by $C_9$ and/or $C_{10}$. In addition to (42), we thus expect\footnote{We define the universality ratios as $R(Y_s)_{\lambda_s^a} = \Gamma(Y_s \to X_\mu \mu)/\Gamma(Y_s \to X_e e)$, assuming a region in $m_{\ell_f}$ below the charmion resonances and sufficiently above the di-muon threshold ($m_{\ell_f} \gtrsim 3$ GeV).}

$$R_{K} \approx R_{K^*} \approx R_{\Lambda_b} \approx R_{(\Lambda_b)_{\mu}} \approx R_{(\Lambda_b)_{e}} \approx \cdots \approx R_K.$$

(43)

Current experimental data hint to sizable NP effects in $R_K$ and $R_{K^*}$ [7-11], consistent with $R_K \approx R_{K^*}$.

Assuming the NP effect to be the same in $R_K$ and $R_{K^*}$, the combined measurements imply $R_{K^{(*)}}^{\text{exp}} = 0.80 \pm 0.05$ (corresponding to $\Delta C_9^{\mu*} = -0.43 \pm 0.11$). This numerical value provides an important constraint on the size of the lepton spurious ($\lambda_q^\mu$): since $\Delta q^{\mu*} = 0(\lambda_q^\mu)^2$, setting $\lambda_q^\mu = 0(10^{-1})$ and $C_V = 0(10^{-2})$, as suggested by the $R_{D^{(*)}}$ fit, the value of $R_{K^{(*)}}$ implies $\lambda_q^\mu = 0(10^{-1})$. 

\[\text{Fig. 3. Predictions for } B(\bar{B}_s \to \tau\bar{\nu}), B(\bar{B}_s \to \pi\tau\bar{\nu}) \text{ and } B(\bar{B}_s \to \tau\pi\bar{\nu}), \text{ all normalized to the corresponding SM expectations, as a function of } \Delta R_{D} - \Delta R_{D^*}. \text{ In gray, the experimental value of } B(\bar{B}_s \to \tau\bar{\nu}) \text{ at } 1\sigma \text{ and } 2\sigma.\]
Within the UV leptoquark completion, the fact that $\Delta R_{K^{(*)}} \equiv R_{K^{(*)}} - 1 < 0$ allows us also to determine a non-trivial relative sign among the $U(2)^3$ breaking terms: according to (20) one has $C_V > 0$, which implies $\Delta q_{\mu}^2 > 0$. Other data involving $b \to s \mu \bar{\mu}$ transitions, such as the measurements of $P_{B_d}^{\mu}$ [63–66] and other differential distributions, also deviate from the SM predictions consistently with $R_{K^{(*)}}$, further supporting the hypothesis of $\lambda_+^V, \lambda_-^V = 0(10^{-1})$. This coincidence in size for quark and lepton spurious points to the interesting possibility of a common origin for the two leading $U(2)^3$ breaking terms.

$b_s \to \mu \bar{\mu}$. Among $b \to s \mu \bar{\mu}$ transitions, a special role is played by $b_s \to \mu \bar{\mu}$, where the chiral enhancement of the scalar amplitude allows us to probe the helicity structure of the NP interaction. The branching ratio normalized to the SM value reads

$$\frac{B(B_s \to \mu \bar{\mu})_{\text{SM}}}{B(B_s \to \mu \bar{\mu})_{\text{SM}}} = 1 + \frac{\Delta c_{10}^2}{C_{10}^2} \chi_s \eta_S \frac{m_t}{m_\mu} \frac{C_{\mu}}{C_V} + \left(1 - \frac{4m_\mu^2}{m_t^2} \right) \chi_s \eta_S \frac{m_t}{m_\mu} \frac{C_{\mu}}{C_V}$$

expressing the deviations in the Wilson coefficients in terms of $\Delta R_{K^{(*)}}$, by means of (41) and (42), leads to

$$\frac{B(B_s \to \mu \bar{\mu})}{B(B_s \to \mu \bar{\mu})_{\text{SM}}} = 1 - \frac{\Delta R_{K^{(*)}}}{0.47 \times 10^{-2}} \left(1 - \chi_s \eta_S \frac{m_t}{m_\mu} \frac{C_{\mu}}{C_V} \right) + \left(1 - \frac{4m_\mu^2}{m_t^2} \right) \chi_s \eta_S \frac{m_t}{m_\mu} \frac{C_{\mu}}{C_V} \left(1 + \frac{\Delta c_{10}^2}{C_{10}^2} \right) \frac{1}{\chi_s \eta_S} \frac{m_t}{m_\mu} \frac{C_{\mu}}{C_V}$$

Current experimental measurements [67–71] yield

$$B(B_s \to \mu \bar{\mu})_{\text{exp}} = 2.72(34) \times 10^{-9},$$

which is about 2.6$\sigma$ below the SM expectation: $B(B_s \to \mu \bar{\mu})_{\text{SM}} = 3.66(14) \times 10^{-9}$ [72]. In Fig. 4, we show the predictions for this observable as a function of $\Delta R_{K^{(*)}}$, for $s_t = 0$ (purple band), and for $s_t = -0.1 \lambda_+^V$ setting $C_5/C_V = 2$ (green band). As can be seen, the deviations in $R_{K^{(*)}}$ are well compatible with the current experimental values of $B(B_s \to \mu \bar{\mu})$ and, if $C_5/C_V$ is large, small values of $s_t$ are favored.

$b \to s t \bar{\mu}$. As far as LFV processes are concerned, the most relevant observable is

$$B(B_s \to \tau^- \mu^+) \approx \frac{\tau_B m_{B_s} f_{B_s}^2 \Gamma_{\tau}^2}{8\pi} \left(1 - \frac{m_s^2}{m_\mu^2} \right) \frac{C_\tau}{C_V} \chi_s \eta_S \frac{m_t}{m_\mu} \frac{C_{\mu}}{C_V} \left(1 + \frac{\Delta c_{10}^2}{C_{10}^2} \right) \frac{1}{\chi_s \eta_S} \frac{m_t}{m_\mu} \frac{C_{\mu}}{C_V}$$

as in $B(B_s \to \tau \bar{\tau})$, the large chiral enhancement of the scalar contribution make it an excellent probe of the helicity structure of the NP effects. Moreover, this observable provides a direct probe of $\Delta q_{\mu}^2$. Setting $C_5/C_V = 0(10^{-2})$ and $\Delta q_{\mu}^2 = 0(10^{-2})$, we find $B(B_s \to \tau^- \mu^+) \approx 2 \times 10^{-6}$, while for $C_5 = 0$ and the same NP parameters, the expected branching fraction is about one order of magnitude smaller. The current experimental limit, $B(B_s \to \tau^- \mu^+) < 4.2 \times 10^{-5}$ (95% CL) [73], is close to the NP predictions when $C_5$ is sizable. Future improvements in this observable will therefore provide very significant constraints.

4.4. $b \to d \ell \bar{\ell}$ and other FCNCs

A key prediction of the minimally broken $U(2)^5$ framework is that NP effects in $b \to s \ell \bar{\ell}$ and $b \to d \ell \bar{\ell}$ transitions scale according to the corresponding CKM factors. More precisely, defining the effective hamiltonian of the leading $b \to d$ FCNC operators as

$$\mathcal{H}_{\text{WET}}^{b \to d} = - \frac{4G_F}{\sqrt{2}} \frac{\alpha}{4\pi} \tilde{V}_{td} V_{td}^* \sum_{i=9,10,5,6} \tilde{c}_i^\ell \tilde{O}_i^\ell,$$

where $\tilde{O}_i^\ell = O_i^\ell [s \to d]$, then it is easy to check that, because of (17),

$$\Delta q_{i=10}^2 = \Delta C_{i=10}^2 \quad \text{and} \quad \tilde{c}_{i=9}^\ell = C_{i=9}^\ell \quad \text{and} \quad \tilde{c}_{i=5}^\ell = C_{i=5}^\ell.$$

These relations lead to a series of accurate predictions which could be tested in various $b \to d \ell \bar{\ell}$ observables.

One of the cleanest test is obtained by means of $B \to \pi \mu \bar{\mu}(e\bar{e})$ decays, where we expect

$$\frac{B(B \to \pi \nu \mu \bar{\nu})_{\text{exp}}}{B(B \to \pi \nu \mu \bar{\nu})_{\text{SM}}} \approx R_{K^{(*)}},$$

where $\Delta q_{\text{pert}}^2$ denotes an interval in $q^2 = m_{\ell \ell}^2$, where perturbative contributions are dominant. The SM prediction for the rate is $B(B^+ \to \pi^+ \mu^+ \bar{\nu})_{\text{SM}} = 1.31(25) \times 10^{-9}$ [74] and $B(B^{+} \to \pi^{+} \mu^{+} \bar{\nu})_{\text{SM}} = 0.72(7) \times 10^{-9}$ [75], to be compared with the LHCb results $B(B^+ \to \pi^+ \mu^+ \bar{\nu})_{16} = 0.91(21) \times 10^{-9}$ and $B(B^+ \to \pi^+ \mu^+ \bar{\nu})_{15,22} = 0.47(11) \times 10^{-9}$ [76]. These measurements deviate from the SM predictions by 1.2$\sigma$ and 2$\sigma$, respectively, and are well consistent with (50) (assuming NP effect in the $e\bar{e}$ mode to be valid).

6 The $q^2$ regions where perturbative contributions dominates over charmonia or light resonance terms are the low-$q^2$ region ($2 \text{ GeV}^2 \leq q^2 \leq 6 \text{ GeV}^2$) and the high-$q^2$ region ($q^2 \gtrsim 15 \text{ GeV}^2$).
Table 1

| Process | Spurion comb. for $dS \to \ell\ell'$ | Bound on $C_V$ |
|---------|-----------------------------------|----------------|
| $K^0 \to \mu\mu$ | $\Delta S_{\mu}^{(\mu)}$ | $C_V \lesssim 0.3$ |
| $K^0 \to e\bar{e}$ | $\Delta S_{e}^{(e)}$ | $C_V \lesssim 0.1(0.2)$ |

subleading), although they are still affected by large errors. Similarly, our framework predicts

$$B(B_s \to \mu\mu) \approx B(B_d \to \mu\mu)_{SM}$$

(51)

Leaving aside $B$ decays, the effective Lagrangian (12) necessarily imply non-standard effects also in $K$ and $\tau$ semi-leptonic decays. Since NP effects in these processes are strongly constrained, it is important to check if these constraints limit the parameter space of the EFT. As far as $\tau$ decays are concerned, the most stringent constraint is obtained by the non-observation of $\tau \to \mu \gamma$. Only the chiral-enhanced contribution to $\tau \to \mu \gamma$, proportional to $C_5$, can be reliably estimated in the EFT, yielding\footnote{The full branching fraction, including both vector and non-chiral-enhanced scalar contributions, was computed in [20] in a specific $U_1$ UV-completion. There, it was found that the additional contributions are much smaller than the ones quoted here if $C_5 \sim C_7$.}

$$B(\tau \to \mu \gamma) \approx \frac{1}{\Gamma_\tau} \frac{\alpha\pi}{256\pi^4} \frac{m_\tau^3 m_\mu^2}{v^4} |C_5 \lambda_\mu^\mu \lambda_\mu^\gamma|^2. \quad (52)$$

Taking $C_5 = O(10^{-2})$ and $\lambda_\mu^\mu = O(10^{-1})$, we find $B(\tau \to \mu \gamma) = \text{few} \times 10^{-9}$, which is below current bounds but within the expected Belle II sensitivity [77].

Finally, the constraints obtained from $K$ decays do not yield significant bounds to our framework. As in the $b \to s \nu \bar{\nu}$ case, NP effects in $s \to d \nu \bar{\nu}$ transitions are forbidden at tree-level if we take $C_{VL} = C_{V}$. On the other hand, $K_1 \to \ell\ell'$ decays receive strong spurion suppressions, resulting in bounds on $C_7$ that are significantly above the preferred values. These are shown in Table 1, together with the parametric spurion dependence of the corresponding $dS \to \ell\ell'$ transition. For comparison, we stress that the preferred values of $C_7$ emerging from the $R_{D^{(*)}}$ fit in Fig. 1, assuming $\lambda_\mu^\mu = 0.1$, is $C_7 = C_{V}/(1 - \lambda_\mu^\nu V_{\mu\tau}^* / V_{\mu\tau}) \sim (1/2) \times 10^{-2}$.

4.5. Charged-current transitions to light leptons

The $U(2)^5$ breaking in the lepton sector could in principle be tested also in charged-current decays to light leptons, both in $b \to c$ and $b \to u$ transitions. The most relevant observables in each category are respectively $R_{D^{(*)}}^{\mu/e} \equiv B(\bar{B} \to D^{(*)}\mu\bar{\nu})/B(\bar{B} \to D^{(*)}\nu\bar{\nu})$ and $B(\bar{B} \to \mu\bar{\nu})$, whose experimental measurements will be improved at Belle II [77]. In contrast to $R_{D^{(*)}}$, scalar contributions to $R_{D^{(*)}}^{\mu/e}$ are extremely suppressed due to the $U(2)^5$ flavor symmetry [see (16)].

The $\mu/e$ LFU ratios can be expressed as

$$\frac{R_{D^{(*)}}^{\mu}}{R_{D^{(*)}}^{e}} \approx |1 + \lambda_\mu^{\mu\mu} e^{C_{V}(\mu)}|^2 + |\lambda_\mu^{\mu\mu} C_{V}(\mu)|^2, \quad (53)$$

where the first term corresponds to the mode with $\nu = \nu_{\tau}$ and the second one with $\nu = \nu_{e}$ and where, similarly to $C_{V}$, we have defined

$$C_{V}^{\mu} \equiv \lambda_\mu^{\mu} C_{V} \left[ 1 - \frac{\Delta S_{\mu}^{(\mu)}}{\lambda_\mu^{\mu}} \frac{V_{\mu\tau}}{V_{\mu\tau}} \right]. \quad (54)$$

Taking $\lambda_\mu^{\mu} = 0(10^{-1})$, $C_{V}^{\mu} = 0(10^{-1})$ and $C_{V}^{\mu} = 0(\lambda_\mu^{\mu} C_{V}) = 0(10^{-2})$, we find that NP corrections to these observables are at most at the per-mille level, hence beyond the near-future experimental sensitivity. This is quite different from what is expected in other NP models addressing the anomalies, such as the scalar leptoquark models considered in [79,80].

It is also worth stressing that the phenomenological condition $\Delta S_{\mu}^{(\mu)} < 0$, required to accommodate $R_{D^{(*)}}$, yields a partial cancellation between the two terms in $C_{V}^{\mu}$. As a result of this cancellation, NP effects are typically at the sub-per-mille level, hence beyond any realistic sensitivity even in a long-term perspective. Similarly, we find possible NP contributions to $B(\bar{B} \to \mu\bar{\nu})$ to be at or below the per-mille level, very far from the experimental reach.

5. Conclusions

The hints of LFU violations observed in $B$ meson decays have shaken many prejudices about physics beyond the SM, opening new directions in model building. One of the most intriguing possibilities is the existence of a link between the (non-standard) dynamics responsible for these anomalies and that responsible for the fermion mass hierarchies. A specific realization of this idea is the hypothesis that, at low energies, the new dynamics manifest via an EFT controlled by an approximate $U(2)^3$ symmetry, with lepton breaking in specific directions in the $U(2)_\ell$ and $U(2)_d$ subgroups.

In this paper we have explored in generality the consequences of this symmetry and symmetry-breaking hypothesis in (semi)leptonic $B$ decays, trying to avoid making additional dynamical assumptions about the origin of the anomalies. As we have shown, the symmetry hypothesis alone leads to a significant reduction in the number of free parameters of the EFT which, in turn, can be translated into stringent predictions on various low-energy observables. The situation is particularly simple in the case of charged currents, where all relevant processes are controlled by two independent combinations of effective couplings. The latter can be determined for instance from $R_{D}$ and $R_{D^{(*)}}$, leading to a series of unambiguous predictions for $B(\bar{B} \to \ell\ell')$, $B(\bar{B} \to \pi\ell\nu)$, polarizion asymmetries in $\bar{B} \to D^{(*)}\tau\bar{\nu}$, as well as other processes. As shown in Fig. 1, the available data on $B(\bar{B} \to \tau\bar{\nu})$ perfectly support the initial hypothesis.

In neutral currents, additional combinations of effective couplings appear, but also in this case a series of stringent predictions, which are genuine consequences of the symmetry and symmetry-breaking hypothesis alone, can be derived. The two most notable ones are: i) the (approximate) universality of the deviations from 1 in $\mu/e$ ratios in short-distance dominated $b \to s\ell\ell$ transitions, leading to (43), and ii) the SM-like CKM scaling of NP effects in $b \to s$ and $b \to d$ transitions, which leads to the relations (50) and (51). If the significance of the current anomalies will increase, the experimental tests of these relations, which are within the reach of current facilities, will provide an invaluable help in clarifying the origin of this intriguing phenomenon.

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Appendix A. Physical parameters in the $U(2)^3$ spurions

We describe here how to remove the unphysical parameters in the spurions appearing in (4). Let us focus first on the quark sector. In the exact $U(2)^3$ limit, the Lagrangian presents a

$$U(2)_q \times U(2)_u \times U(2)_d \times U(1)_B,$$

(A.1)

flavor symmetry. Since, by assumption, this symmetry is only broken by $\Delta_{u,d}$ and $V_q$, field transformations under this symmetry only modify these spurions. In particular, the $U(2)^3$ transformation

$$q_i \to U_q q_i, \quad u_i \to U_u u_i, \quad d_i \to U_d d_i,$$

(A.2)

with $U_{Q,U,D}$ being general $2 \times 2$ unitary matrices, let us bring $\Delta_{u,d}$ to the following form

$$\Delta_u \to U_Q^\dagger \Delta_u U_Q = O_u \hat{\Delta}_u,$$

$$\Delta_d \to U_D^\dagger \Delta_d U_D = \hat{\Delta}_d.$$  

(A.3)

Here $\hat{\Delta}_u$ and $\hat{\Delta}_d$ are diagonal positive matrices and $O_u$ is an orthogonal matrix. This is the well-known result that in the SM with two families, only the Cabibbo matrix is physical and there are no observable CP-violating phases. A residual $U(1)_R$ symmetry, corresponding to baryon number for light quarks, remains unbroken by these spurions. In terms of degrees of freedom, we started with 8 real parameters and 8 phases in $\Delta_{u,d}$, and we used the free freedom of the $U(2)^3\times U(1)_R$ symmetry to remove 11 parameters. In the end, 5 real parameters remain: 4 quark masses and one mixing angle.

The field redefinitions performed so far also redefine the $V_q$ spurion: $V_q \to \tilde{V}_q = U_Q^\dagger V_q$. We can decompose it as

$$\tilde{V}_q = e^{i\alpha_q} |V_q| U_q \tilde{n}, \quad \tilde{n} = \left( \begin{array}{c} 0 \\ 1 \end{array} \right),$$

(A.4)

with $U_q$ defined as in (8). Without loss of generality, the phase $\alpha_q$ can be absorbed into a redefinition of $x_{l,b}$. Moreover, we can use the $[U_{B_i} \times U(1)_B]/U(1)_B$ symmetry transformation to remove one phase combination from the two phases in $x_{l,b}$. We use this freedom to set equal phases for $x_q$ and $x_{q'}$. Finally, we can absorb the phase of $y_{l,b}$ into a redefinition of the $\xi_k$ and $b_k$ fields. This final redefinition modifies the phase of any $U(2)^2$-preserving NP interaction, but it does not affect the rest of the SM Lagrangian. After all these field redefinitions, we end up with the following quark Yukawa matrices

$$Y_u = |y_t| \left( \begin{array}{cc} O_u^\dagger & U_q |V_q| \phi_q e^{i\delta_q} \tilde{\omega} \\ 0 & 1 \end{array} \right),$$

$$Y_d = |y_b| \left( \begin{array}{cc} \hat{\Delta}_d^\dagger & U_q |V_q| \phi_q e^{i\delta_q} \tilde{\omega} \\ 0 & 1 \end{array} \right).$$

(A.5)

A final $U(2)_q$ transformation, $q_i \to U_q^\dagger q_i$, is sufficient to bring these matrices to the form in (7). The same procedure can be applied to the lepton sector, however in this case there are three differences with respect to the quark case (in the limit of vanishing neutrino masses): i) a subgroup of the $U(2)_e \times U(2)_\tau$ symmetry is enough to make $\Delta_e$ diagonal and positive, with no extra orthogonal matrices; ii) the remaining freedom let us reduce $U_e$ to an orthogonal matrix, which we denote as $O_e$; iii) the $[U(1)_L \times U(1)_L]/U(1)_L$ transformation can be used to remove the phase in $x_{e,\tau}$, leaving no free phases. After an appropriate $U(2)_L$ transformation, it is straightforward to arrive to the form in (7).

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