Quark flavour mixing with right-handed currents: an effective theory approach

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Abstract

The impact of right-handed currents in both charged- and neutral-current flavour-violating processes is analysed by means of an effective theory approach. More explicitly, we analyse the structure of dimension-six operators assuming a left-right symmetric flavour group, commuting with an underlying $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ global symmetry, broken only by two Yukawa couplings. The model contains a new unitary matrix controlling flavour-mixing in the right-handed sector. We determine the structure of this matrix by charged-current data, where the tension between inclusive and exclusive determinations of $|V_{ub}|$ can be solved. Having determined the size and the flavour structure of right-handed currents, we investigate how they would manifest themselves in neutral current processes, including particle-antiparticle mixing, $Z \to b\bar{b}$, $B_{s,d} \to \mu^+\mu^-$, $B \to \{X_s, K, K^*\} \nu\bar{\nu}$, and $K \to \pi\nu\bar{\nu}$ decays. The possibility to explain a non-standard CP-violating phase in $B_s$ mixing in this context, and the comparison with other predictive new-physics frameworks addressing the same problem, is also discussed. While a large $S_{\psi\phi}$ asymmetry can easily be accommodated, we point out a tension in this framework between $|V_{ub}|$ and $S_{\psi K}$.

1 Introduction

One of the main properties of the Standard Model (SM) regarding flavour violating processes is the left-handed structure of the charged currents that is in accordance with the maximal violation of parity observed in low energy processes. Left-handed charged currents encode at the level of the Lagrangian the full information about flavour mixing and CP violation represented compactly by the CKM matrix. Due to the GIM mechanism this structure has automatically profound implications for the pattern of FCNC processes that seems to be remarkably in accordance with the present data within theoretical and experimental uncertainties, bearing in mind certain anomalies which will be discussed below.

Yet, the SM is expected to be only the low-energy limit of a more fundamental theory in which parity could be a good symmetry implying the existence of right-handed charged
currents. Prominent examples of such fundamental theories are left-right symmetric models on which a rich literature exists.

Left-right symmetric models were born 35 years ago and extensive analyses of many observables can be found in the literature (see e.g., and references therein). Renewed theoretical interest in models with an underlying SU(2)_L × SU(2)_R global symmetry has also been motivated by Higgsless models. However, the recent phenomenological interest in making another look at the right-handed currents in general, and not necessarily in the context of a given left-right symmetric model, originated in tensions between inclusive and exclusive determinations of the elements of the CKM matrix |V_{ub}| and |V_{cb}|. It could be that these tensions are due to the underestimate of theoretical and/or experimental uncertainties. Yet, it is a fact, as pointed out and analyzed recently in particular in Ref. [12, 13], that the presence of right-handed currents could either remove or significantly weaken some of these tensions, especially in the case of |V_{ub}|.

Assuming that right-handed currents provide the solution to the problem at hand, there is an important question whether the strength of the right-handed currents required for this purpose is consistent with other observables and whether it implies new effects somewhere else that could be used to test this idea more globally.

In the present paper we make still another look at the effects of the right-handed currents in low energy processes, without specifying the fundamental theory in detail but only assuming its global symmetry and the pattern of its breakdown. Specifically we address the following questions:

- What is the allowed structure of the right-handed matrix that governs flavour violating processes in the right-handed sector and reduces the tension observed in the inclusive and exclusive determinations of |V_{ub}|?
- What is the impact of right-handed currents, in combination with SM left-handed currents, on particle-antiparticle mixing?
- Can in this context the known Z → b̅b “anomaly” be solved?
- Can such effects be seen in rare FCNC decays such as B_{s,d} → μ^+μ^−, B → \{X_s, K, K^*\}ν̅ν̅ and K → πν̅ν̅?

As there have been already numerous analyses of right-handed currents in the literature, it is mandatory for us to state what is new in our paper:

- In the spirit of [14] we use an effective theory approach to describe the effects of right-handed currents in flavour violating processes. In fact our work could be considered to be a simple generalization of the usual MFV framework to include right-handed currents. Indeed the main model-dependent assumption of our analysis is the hypothesis that the left-right symmetric flavour group is broken only by two Yukawa couplings.
- We determine the new unitary matrix ˜V that controls flavour-mixing in the right-handed sector by using the data on the tree level charged current transitions u → d, u → s, b → u and b → c and its unitarity. Here the novel feature of our analysis, as compared to [12, 15], is the determination of the full right-handed matrix and not only of its b → c, u elements.
- We point out that the elements of this matrix can be further constrained through the FCNC processes, that we analyse in detail. Here our minimalistic assumption about
the breaking of the flavour symmetry by only two Yukawa couplings plays a key role, and distinguishes our work from most of the existing analyses.

- We point out that the explanation of the different values of $|V_{ub}|$ from inclusive and exclusive semi-leptonic decays with the help of right-handed currents, implying large value of $|V_{ub}|$, strengthens the tension (already existing in the SM) between $\sin 2\beta$ and $S_{\psi K_S}$. This tension cannot be solved through the new CP-violating effects in the right-handed matrix as the contributions of right-handed currents to $B^0_d - \bar{B}^0_d$ mixing turn out to be strongly suppressed due to the $\varepsilon_K$ constraint and the desire to explain the enhanced value of the $S_{\psi K}$ asymmetry in the $B_s$ system.

Our paper is organized as follows. In Section 2 we present the electroweak symmetry and the field content of our effective theory. Here we also introduce the right-handed (RH) matrix $\tilde{V}$ and discuss some of its properties. In Section 3 we make a closer look at the RH charged currents. In Section 4 we use the present knowledge on $u \to d$, $s \to d$, $b \to c$ and $b \to u$ transitions to obtain upper bounds on most of the elements of the matrix $\tilde{V}$. Section 5 is devoted to a general discussion of dimension-six operators relevant to neutral currents. The impact of these operators is then analysed in particle-antiparticle mixing (Section 6) and in processes sensitive to $Z$-mediated neutral currents (Section 7). A comparison of the results obtained in the present framework with those obtained under the MFV hypothesis, and in explicit left-right models, is presented in Section 8. We summarize our main results and conclude in Section 9.

2 The model

2.1 Electroweak symmetry and field content

The starting point of our analysis is the assumption that the SM is the low-energy limit of a more fundamental theory. We don’t know the exact structure of this theory, but we assume that in the high-energy limit it is left-right symmetric. The difference of left-handed (LH) and right-handed sectors observed in the SM is only a low-energy property, due to appropriate symmetry-breaking terms.

In particular, we assume that the theory has a $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ global symmetry, explicitly broken only in the Yukawa sector and by the $U(1)_Y$ gauge coupling. Under this symmetry the SM quark fields can be grouped into three sets of LH or RH doublets with $B-L$ charge 1/3:

$$Q^i_L = \begin{pmatrix} u^i_L \\ d^i_L \end{pmatrix}, \quad Q^i_R = \begin{pmatrix} u^i_R \\ d^i_R \end{pmatrix}, \quad i = 1 \ldots 3. \quad (1)$$

With this assignment the SM hypercharge is given by $Y = T^3_{3R} + (B - L)/2$. In order to recover the SM electroweak gauge group, we assume that only the $SU(2)_L$ and $U(1)_Y$ subgroups of $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ are effectively gauged below the TeV scale. In close analogy we can introduce three sets of LH and RH leptons, $L^i_L$ and $L^i_R$ (including three RH neutrinos), with $B-L = -1$.

The breaking of the electroweak symmetry is achieved via the spontaneous breaking of $SU(2)_L \times SU(2)_R$ into the vectorial subgroup $SU(2)_{L+R}$ at the electroweak scale. For
simplicity, we provide an explicit description of this breaking introducing a SM-like Higgs field transforming as \((2,\bar{2})\) of \(SU(2)_L \times SU(2)_R\), with non-vanishing vacuum expectations value (vev):

\[
H \rightarrow U_L H U_R^\dagger, \quad U_{L(R)} \in SU(2)_{L(R)} , \quad \langle H \rangle = v \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix} .
\] (2)

Introducing the kinetic term

\[
\mathcal{L}_{\text{kin}}^{\text{Higgs}} = \frac{1}{4} \text{Tr}[(D_\mu H)^\dagger D_\mu H] ,
\] (3)

where \(D_\mu H = \partial_\mu H - igW_\mu^a T^a H + ig' H T^3 B_\mu\), we recover the standard tree-level expressions of \(W\) and \(Z\) masses for \(v \approx 246\) GeV. However, most of the following discussion on flavour-violating effective operators applies as well to models where the \(SU(2)_L \times SU(2)_R\) breaking is achieved via a more complicated Higgs sector, or even without a fundamental Higgs field. In the latter case \(H\) is replaced by \(v \times U\), where \(U\) is the unitary field, transforming as \((2,\bar{2})\) of \(SU(2)_L \times SU(2)_R\), that encodes the three Goldstone bosons (see e.g. Ref. [16]).

### 2.2 The quark flavour symmetry

As far as the quark flavour structure is concerned, we assume an \(SU(3)_L \times SU(3)_R\) flavour symmetry, with the fields transforming as

\[
Q_{L(R)} \rightarrow f_{L(R)} Q_{L(R)} , \quad f_{L(R)} \in SU(3)_{L(R)} .
\] (4)

In order to generate different masses for up- and down-type quarks we introduce two spurion fields, \(P_u(d)\), transforming as \(U_R P_{u(d)} U_R^\dagger\) under \(SU(2)_R\), whose background values break the custodial \(SU(2)_{L+R}\) symmetry:

\[
P_u = \begin{pmatrix} 1 & 0 \\ 0 & 0 \end{pmatrix} , \quad P_d = \begin{pmatrix} 0 & 0 \\ 0 & 1 \end{pmatrix} .
\] (5)

In addition, we assume that the flavour symmetry is broken by two spurions, both transforming as \((3,\bar{3})\) under \(SU(3)_L \times SU(3)_R\):

\[
Y_{u(d)} \rightarrow Y_{u(d)} f_R^\dagger .
\] (6)

Finally, we assume an additional \(U(1)\) symmetry, under which \(Y_u\) and \(Y_d\) have different charges, and \(P_u\) and \(P_d\) have the corresponding opposite charges. This symmetry structure implies the following invariant quark Yukawa interactions:

\[
\mathcal{L}_Y = \bar{Q}_L H Y_u P_u Q_R + \bar{Q}_L H Y_d P_d Q_R + \text{h.c.}
\] (7)

This is equivalent to the SM Lagrangian when we take into account the structure of \(P_u\) and \(P_d\). In models with two Higgs doublets the two spurions \(P_u\) and \(P_d\) can be interpreted as the remnant of the different vevs of the two Higgs fields. Assuming heavy masses and similar vevs for these Higgs fields, this picture cannot be distinguished from the one-Higgs doublet case in our effective-theory approach.
Rotating $Q_L$ and $Q_R$ in flavour space we can always choose a quark basis where one of the two Yukawa couplings is diagonal. We can also rotate the relative phases of the quark fields to make this diagonal matrix real. Choosing the basis where $Y_d$ is diagonal we can write

$$Y_d \mid_{d\text{-base}} = \lambda_d , \quad \lambda_d = \frac{\sqrt{2}}{v} \text{diag}(m_d, m_s, m_b) \equiv \text{diag}(y_d, y_s, y_b) ,$$

$$Y_u \mid_{d\text{-base}} = V^\dagger \lambda_u \tilde{V} , \quad \lambda_u = \frac{\sqrt{2}}{v} \text{diag}(m_u, m_c, m_t) \equiv \text{diag}(y_u, y_c, y_t) ,$$

where $V$ and $\tilde{V}$ are two unitary complex $3 \times 3$ mixing matrices. $V$ can be identified with the CKM matrix, while $\tilde{V}$ is a new unitary mixing matrix that parametrizes the misalignment of $Y_u$ and $Y_d$ in the RH sector. In such a basis, compatible with the $SU(3)_L \times SU(3)_R$ flavour symmetry, the mass terms generated by $L Y$ once the Higgs gets a vev are:

$$L_y \mid_{d\text{-base}} = v \bar{u} L V^\dagger \lambda_u \tilde{V} u R + v \bar{d} \lambda_d d R + \text{h.c.}$$

In order to diagonalize the mass terms for the up-quarks we need to perform the following (flavour-breaking) rotations

$$u_L \rightarrow u'_L = V u_L , \quad u_R \rightarrow u'_R = \tilde{V} u_R ,$$

with $u'_{L,R}$ denoting the mass-eigenstate fields.

### 2.3 RH mixing matrix and FCNC spurions

The new mixing matrix $\tilde{V}$ can be parametrized in terms of 3 real mixing angles and 6 complex phases. Adopting the standard CKM phase convention, where the 5 relative phases of the quark fields are adjusted to remove 5 complex phases from the CKM matrix, we have no more freedom to remove the 6 complex phases from $\tilde{V}$. In the standard CKM basis $\tilde{V}$ can be parametrized as follows

$$\tilde{V} = D_U \tilde{V}_0 D_D^T ,$$

where $\tilde{V}_0$ is a “CKM-like” mixing matrix, containing only one non-trivial phase and $D_{U,D}$ are diagonal matrices containing the remaining CP-violating phases. For reasons that will become clear in the following, it is convenient to attribute the non-trivial phase of $\tilde{V}_0$ to the 2–3 mixing, such that

$$\tilde{V}_0 = \begin{pmatrix} \tilde{c}_{12} \tilde{c}_{13} & \tilde{s}_{12} \tilde{c}_{13} & \tilde{s}_{13} \\ -\tilde{s}_{12} \tilde{c}_{23} - \tilde{c}_{12} \tilde{s}_{23} \tilde{s}_{13} e^{-i\phi} & \tilde{c}_{12} \tilde{c}_{23} - \tilde{s}_{12} \tilde{s}_{23} \tilde{s}_{13} e^{-i\phi} & \tilde{s}_{23} \tilde{c}_{13} e^{-i\phi} \\ -\tilde{c}_{12} \tilde{c}_{23} \tilde{s}_{13} + \tilde{s}_{12} \tilde{s}_{23} e^{i\phi} & -\tilde{c}_{12} \tilde{s}_{23} \tilde{s}_{13} - \tilde{s}_{12} \tilde{c}_{23} e^{i\phi} & \tilde{c}_{23} \tilde{c}_{13} \end{pmatrix} ,$$

and

$$D_U = \text{diag}(1, e^{i\phi_2}, e^{i\phi_3}) , \quad D_D = \text{diag}(e^{i\phi_4}, e^{i\phi_5}, e^{i\phi_6}) .$$

Combining the basic spurions $Y_u$ and $Y_d$ we can build symmetry-breaking terms transforming as $(8, 1)$ or $(1, 8)$ under the $SU(3)_L \times SU(3)_R$ flavour group. These spurions control the strength of FCNCs in the model. Since the only large Yukawa coupling is the top-one,
only two of such terms are phenomenologically relevant: \( Y_u Y_u^\dagger \sim (8, 1) \) and \( Y_u^\dagger Y_u \sim (1, 8) \). Their explicit expressions in the mass-eigenstate basis of down quarks are

\[
(Y_u Y_u^\dagger)_{i \neq j \mid d\text{-base}} = (V^\dagger \lambda_2^0 V)_{ij} \approx y_i^2 V_{3i}^* V_{3j},
\]

(14)

\[
(Y_u^\dagger Y_u)_{i \neq j \mid d\text{-base}} = (\tilde{V}^\dagger \lambda_2^0 \tilde{V})_{ij} \approx y_i^2 e^{i(\phi_d^i - \phi_d^j)} (\tilde{V}_0)_3^i (\tilde{V}_0)_3^j.
\]

(15)

The \( Y_u Y_u^\dagger \) term, which appears in LH mediated FCNCs, has exactly the same structure as in the MFV framework \[14\]. The \( Y_u^\dagger Y_u \) term is a new spurion characterizing the strength of RH mediated FCNCs in this model. To make contact with the analysis of Ref. \[17\], where the MFV flavour group \((SU(3)_{QL} \times SU(3)_{uR} \times SU(3)_{dR})\) with non-minimal breaking terms has been considered, our framework corresponds to the introduction of a single non-MFV spurion, \( \tilde{V} \), that transforms as \((\bar{3}, \bar{3})\) under \( SU(3)_{uR} \times SU(3)_{dR} \) and it is constrained to be a unitary matrix.

### 3 A first look at the dimension-six operators

#### 3.1 Preliminaries

As anticipated, we do not specify the ultraviolet completion of the model. We proceed encoding the effects of the high-energy degrees of freedom by means of the effective Lagrangian

\[ L_{\text{eff}} = L_{\text{SM}} + \frac{1}{\Lambda^2} \sum c_i O_i^{(6)}, \]

(16)

where the \( O_i^{(6)} \) are dimension-six effective operators compatible with the symmetries discussed before. Here \( \Lambda \) is an effective scale, expected to be of \( O(1 \text{ TeV}) \), and the \( c_i \) are dimensionless coefficients.

In order to build the basis of relevant effective operators, it is first convenient to look at the quark bilinear currents invariant under the flavour symmetry defined above. Introducing terms with at most two Yukawa spurions (with no more than one \( Y_d \)), and denoting with \( \Gamma \) a generic Dirac structure, we have

\[
O(Y^0) : \quad \bar{Q}_L \Gamma Q_L, \quad \bar{Q}_R \Gamma Q_R, \quad (17)
\]

\[
O(Y^1) : \quad \bar{Q}_L \Gamma Y_u P_u Q_R, \quad \bar{Q}_L \Gamma Y_d P_d Q_R, \quad (18)
\]

\[
O(Y^2) : \quad \bar{Q}_L \Gamma Y_u^\dagger Y_u Q_L, \quad \bar{Q}_R \Gamma Y_u^\dagger Y_u Q_R. \quad (19)
\]

Most of these bilinear structures are allowed also in the MFV case. The only two notable exceptions are: i) the charged-current component of the RH bilinear \( \bar{Q}_R \Gamma Q_R \), and ii) charged- and neutral-current components of \( \bar{Q}_R \Gamma Y_u^\dagger Y_u Q_R \). These two will play the central role in our paper. We first discuss charged currents, postponing the analysis of neutral currents to Section 5.

#### 3.2 Modification of charged currents

Our first goal is to understand which operators can probe the rotation in the RH sector, namely the flavour-mixing matrix \( \tilde{V} \) that appears in the charged-current component of the
bilinear $\bar{Q}_R \Gamma Q_R$. If we consider operators with only two quark fields, and we ignore RH neutrinos (assuming they are heavy), the list of relevant operators is quite small:

$$O^{(6)}_{R_{\ell 1}} = \bar{Q}_R \gamma^\mu \tau^i \bar{L}_L \gamma_\mu \tau^i L_L,$$
$$O^{(6)}_{R_{h 1}} = i \bar{Q}_R \gamma^\mu H^\dagger D_\mu H Q_R , \quad O^{(6)}_{R_{h 2}} = i \bar{Q}_R \gamma^\mu \gamma_\mu \tau^i \bar{Q}_R \text{Tr} \left( H^\dagger D_\mu H \right).$$

Most important, all these operators are equivalent as far as the quark-lepton charged-current interactions are concerned. In the case of $O^{(6)}_{R_{h 1}}$ we generate an effective coupling of the RH quark current to the $W$ field after the breaking of the electroweak symmetry: integrating out the $W$ leads to a quark-lepton charged-current interaction identical to the one of $O^{(6)}_{R_{\ell 1}}$.

The resulting effective quark-lepton charged-current interaction obtained integrating out the $W$ at the tree-level can be written as

$$L_{\text{c.c.}}^{\text{eff}} = - \frac{g^2}{2 M_W^2} + \frac{c_L}{\Lambda^2} \bar{u}_L \gamma^\mu d_L \bar{\ell}_L \gamma_\mu \nu_L + \frac{c_R}{\Lambda^2} \bar{u}_R \gamma^\mu \bar{d}_R \bar{\ell}_L \gamma_\mu \nu_L + \text{h.c.}$$

In the limit $c_L = c_R = 0$ we recover the usual SM result. The term proportional to $c_R$ is the result of the new operators in Eq. (20): $c_R = -2(c_{R_{\ell 1}} + 2c_{R_{h 2}} - c_{R_{h 1}})$. For completeness, we have also included a possible modification of the LH interaction, parametrized by $c_L$. This is naturally induced by operators obtained from Eq. (20) with $Q_R \rightarrow Q_L$.

In principle, charged-current interactions are potentially sensitive also to operators written in terms of the bilinears in Eqs. (18)–(19). However, as long as we are interested in processes where the up-type quarks are of the first two generations, these terms are safely negligible, being suppressed by small Yukawa couplings.

Rotating the up-type fields to the mass-eigenstate basis by means of Eq. (10), and omitting the prime indices for simplicity, we can finally write

$$L_{\text{c.c.}}^{\text{eff}} = - \frac{4 G_F}{\sqrt{2}} \bar{u}_L \gamma^\mu \left( (1 + \epsilon_L) V P_L + \epsilon_R \tilde{V} P_R \right) d_L \left( \bar{\ell}_L \gamma_\mu \nu_L \right) + \text{h.c.}$$

where

$$P_L = \frac{1 - \gamma_5}{2}, \quad P_R = \frac{1 + \gamma_5}{2},$$
$$\epsilon_R = - \frac{c_R v^2}{2 \Lambda^2} = \frac{v^2}{\Lambda^2} \left( c_{R_{\ell 1}} + 2c_{R_{h 2}} - c_{R_{h 1}} \right), \quad \epsilon_L = - \frac{c_L v^2}{2 \Lambda^2}.$$
4.1 Bounds from $u \rightarrow d$ and $s \rightarrow d$ transitions

Within the SM the leading constraints on $|V_{ud}|$ are derived from super-allowed $(0^+ \rightarrow 0^+)$ nuclear beta decays and by the pion decay ($\pi \rightarrow e\nu$) [18]:

$$|V_{ud}(0^+ \rightarrow 0^+)|_{\text{exp}}^{\text{SM}} = 0.97425(022) ,$$

$$|V_{ud}(\pi \rightarrow e\nu)|_{\text{exp}}^{\text{SM}} = 0.97410(260) .$$

By construction, the super-allowed nuclear beta decays are sensitive only to the $u \rightarrow d$ vector current, while the pion decay is sensitive only the $u \rightarrow d$ axial current. As a result, the corresponding constraints can be implemented in our effective theory via the conditions

$$|1 + \epsilon_L V_{ud} + \epsilon_R \tilde{V}_{ud}| = |V_{ud}(0^+ \rightarrow 0^+)|_{\text{exp}}^{\text{SM}} ,$$

$$|1 + \epsilon_L V_{ud} - \epsilon_R \tilde{V}_{ud}| = |V_{ud}(\pi \rightarrow e\nu)|_{\text{exp}}^{\text{SM}} .$$

In principle there is also a constraint from the neutron beta decay. However, the situation of the neutron life-time measurements is quite confusing at present, and this constraint does not add an additional significant new information.

Since we expect $\epsilon_{L,R} \ll 1$, we can expand the above equations to first order in $\epsilon_{L,R}$. For simplicity, we also assume $\epsilon_{L,R}$ to be real. Solving the above constraints under these conditions leads to

$$|(1 + \epsilon_L)V_{ud}| = 0.9742 \pm 0.0013 , \quad \epsilon_R \text{ Re} \left( \frac{\tilde{V}_{ud}}{V_{ud}} \right) = (0.1 \pm 1.3) \times 10^{-3} .$$

The leading constraints on the vector and the axial $s \rightarrow u$ currents are derived from $K \rightarrow \pi\ell\nu$ and $K \rightarrow \mu\nu$ decays, respectively. Using the SM results obtained in [19],

$$|V_{us}(K \rightarrow \pi\ell\nu)|_{\text{exp}}^{\text{SM}} = 0.2243(12) ,$$

$$|V_{us}(K \rightarrow \mu\nu)|_{\text{exp}}^{\text{SM}} = 0.2252(13) ,$$

we can impose the following conditions

$$|1 + \epsilon_L V_{us} + \epsilon_R \tilde{V}_{us}| = |V_{us}(K \rightarrow \pi\ell\nu)|_{\text{exp}}^{\text{SM}} ,$$

$$|1 + \epsilon_L V_{us} - \epsilon_R \tilde{V}_{us}| = |V_{us}(K \rightarrow \mu\nu)|_{\text{exp}}^{\text{SM}} .$$

Proceeding as in the $u \rightarrow d$ case we find

$$|(1 + \epsilon_L)V_{us}| = 0.2248 \pm 0.0009 , \quad \epsilon_R \text{ Re} \left( \frac{\tilde{V}_{us}}{V_{us}} \right) = -(2.0 \pm 3.9) \times 10^{-3} .$$

Since we know that $|V_{ab}| = O(10^{-3})$, we can neglect $|V_{ab}|^2$ in the CKM unitarity relation $|V_{ud}|^2 + |V_{us}|^2 + |V_{ab}|^2 = 1$, and impose that $|V_{ud}|^2 + |V_{us}|^2 = 1 + O(10^{-4})$. This allows us to obtain a determination of $\epsilon_L$ at the $10^{-3}$ level starting from the constraints on $|(1 + \epsilon_L)V_{ud(s)}|$ in Eq. (29) and (34):

$$\epsilon_L = \left[ (1 + \epsilon_L)^2 (|V_{ud}|^2 + |V_{us}|^2 + |V_{ab}|^2) \right]^{1/2} - 1 = -(0.2 \pm 1.2) \times 10^{-3} .$$
Using this result back in Eq. (29) and (34) we finally obtain:

\[ |V_{ud}| = 0.9742 \pm 0.0013 \, , \quad \epsilon_R \Re(\tilde{V}_{ud}) = (0.1 \pm 1.3) \times 10^{-3} \, , \quad (36) \]

\[ |V_{us}| = 0.2248 \pm 0.0009 \, , \quad \epsilon_R \Re(\tilde{V}_{us}) = -(0.5 \pm 0.9) \times 10^{-3} \, . \quad (37) \]

Note that if \( \epsilon_R = \mathcal{O}(10^{-3}) \), as expected by naïve dimensional analysis for new-physics at the TeV scale, the above results do not imply small mixing angles among the first two generations in the RH sector. Similar conclusions from the analysis of RH currents in semileptonic K decays were reached also in [20]. As we will show in the following, this expectation is confirmed also by our analysis of FCNC processes.

4.2 Bounds from \( b \to c \) and \( b \to u \) transitions

Given the smallness of \( \epsilon_L \) derived in Eq. (35), we can neglect it in the determination of the matrix elements entering \( b \to c \) and \( b \to u \) transitions, where the best experimental errors are at least of \( \mathcal{O}(1\%) \).

First we summarize the bounds from \( b \to c \) transitions, starting with the consideration of the inclusive decay \( B \to X_c \ell \bar{\nu}_\ell \). The total rate is easy to handle as the obtained correction, being dependent on the fraction of the RH mixing matrix element and the CKM matrix element, can be factorized in comparison to the SM tree level decay width. We proceed in the following way: the value for \( V_{cb} \), obtained from the comparison of the SM decay width and the experimental data, must be equivalent to our effective \( V_{cb} \) including the RH contributions. We then obtain

\[ (|V_{cb}|_{\text{SM-exp}})^{\text{incl}} = |V_{cb}|^2 \left[ 1 + |\epsilon_R|^2 \left| \frac{\tilde{V}_{cb}}{V_{cb}} \right|^2 - r_{\text{int}} \Re \left( \epsilon_R \frac{\tilde{V}_{cb}}{V_{cb}} \right) \right] \, , \quad (38) \]

where [21]

\[ |V_{cb}|_{\text{SM-exp}}^{\text{incl}} = (41.54 \pm 0.73) \times 10^{-3} \, , \quad (39) \]

and the strength of the interference \( r_{\text{int}} \) is given by

\[ r_{\text{int}} = 16 \frac{m_c}{m_b} \frac{h(m_b/m_c)}{f(m_c/m_b)} \, , \quad (40) \]

with \( f(x) = 1 - 8x^2 + 8x^6 - x^8 - 24x^4 \log x \) and \( h(x) = 1 - 3x^2 + 3x^4 + x^6 + 6(x^2 + x^4) \log x \). Numerically \( r_{\text{int}} = 0.97 \times 10^{-3} \), thus the impact of the RH current in the inclusive decay turns out to be very small. This is consistent with Ref. [15], where a detailed analysis of the inclusive differential distributions has been performed.

Concerning the exclusive decays \( B \to D^* \ell \bar{\nu}_\ell \) and \( B \to D \ell \bar{\nu}_\ell \), the consideration of the differential decay rate turns out to be more useful as here data has been determined by various experiments. Furthermore the consideration of the heavy-quark limit yields an easy description, where just one form factor, the Isgur-Wise function, has to be taken into account close to the zero-recoil point limit. For completeness, we collect the SM results for the differential decay rates close to this kinematical point

\[ \frac{d\Gamma(B \to D^* \ell \bar{\nu}_\ell)}{dw} = \frac{G_F^2}{48\pi^3} |V_{cb}|^2 m_{D^*}^3 (w^2 - 1)^{1/2} P(w) |\mathcal{F}(w)|^2 \, , \quad (41) \]
\[
\frac{d\Gamma(B \to D\ell\bar{\nu}_\ell)}{dw} = \frac{G_F^2}{4\pi^3} |V_{cb}|^2 (m_B + m_D)^2 m_{D^*}^3 (w^2 - 1)^{3/2} |G(w)|^2 ,
\]

where in the B meson rest frame \( w = E_{D^*}/m_{D^*} \), and the zero-recoil point limit corresponds to \( w = 1 \). Here \( P(w) \) denotes the phase space factor and \( F(w) \) and \( G(w) \) are the hadronic form factors. For \( w = 1 \), when the momentum transfer of the leptons is at its maximum, \( P(1) = 12(m_B - m_{D^*})^2 \). From a fit of the kinematical distribution around \( w = 1 \) the experiments determine with high accuracy the products \( F(1)|V_{cb}| \) and \( G(1)|V_{cb}| \) as well as the curvature of the form factors, obtaining \([21]\).

\[
\begin{align*}
F(1)|V_{cb}|_{SM-exp}^{B \to D^*} &= (35.41 \pm 0.52) \times 10^{-3}, \\
G(1)|V_{cb}|_{SM-exp}^{B \to D^*} &= (42.4 \pm 1.6) \times 10^{-3} .
\end{align*}
\]

In this kinematical limit \( B \to D^*\ell\bar{\nu}_\ell \) and \( B \to D\ell\nu_\ell \) decays involve only axial and vector contributions, respectively. As a result, it is easy to include the RH current contribution. In analogy to the inclusive case, hence our conditions read

\[
\begin{align*}
|V_{cb}|_{SM-exp}^{B \to D^*} &= |V_{cb} - \epsilon_R \tilde{V}_{cb}| , \\
|V_{cb}|_{SM-exp}^{B \to D} &= |V_{cb} + \epsilon_R \tilde{V}_{cb}| .
\end{align*}
\]

In order to implement these constraints we need to specify the values of the form factors at \( w = 1 \). Using the lattice determinations \( G(1) = 1.074 \pm 0.018 \pm 0.0016 \) \([22]\), \( F(1) = 0.921 \pm 0.013 \pm 0.0020 \) \([23]\), leads to

\[
|V_{cb}|_{SM-exp}^{B \to D^*} = (39.4 \pm 1.7) \times 10^{-3} , \quad |V_{cb}|_{SM-exp}^{B \to D} = (38.3 \pm 1.2) \times 10^{-3} .
\]

Performing a global fit to \( V_{cb} \) and \( \epsilon_R \tilde{V}_{cb} \) using the three constraints in Eqs. (39), (45), and (46), we then obtain

\[
|V_{cb}| = (40.7 \pm 0.6) \times 10^{-3} , \quad \epsilon_R \text{ Re} \left( \frac{\tilde{V}_{cb}}{V_{cb}} \right) = (2.5 \pm 2.5) \times 10^{-2} ,
\]

with a modest correlation \( (\rho = 0.16) \). This finally implies

\[
\epsilon_R \text{ Re}(\tilde{V}_{cb}) = (1.0 \pm 1.0) \times 10^{-3} .
\]

In this case the \( \chi^2 \) of the fit is not good \( (\chi^2/N_{dof} = 4.3) \), as also in the SM, because both of the exclusive values in Eq. (47) are below the inclusive one. This result cannot be explained in terms of RH currents. As pointed out in Ref. [24], the inconsistency among the different determinations of \( V_{cb} \) is likely to be due to an overestimate of \( G(1) \) on the Lattice. Lowering the central value to \( G(1) = 0.86 \), as suggested in [24], and keeping the same error, leads to a much better fit \( (\chi^2/N_{dof} = 0.9) \). Since the result for \( \epsilon_R \text{ Re}(\tilde{V}_{cb}) \) obtained in this case is perfectly consistent with the one in [49], in the following we will use Eq. (49) as reference value.

We now proceed analysing the constraints from \( b \to u \) transitions. As far as the inclusive rate is concerned, the structure can be obtained in a straightforward way from the \( b \to c \) case replacing \( c \to u \). Here the interference term is totally negligible, so we obtain the condition

\[
(|V_{ub}|_{SM-exp}^{incl})^2 = (|V_{ub}|^2 + |\epsilon_R|^2|\tilde{V}_{ub}|^2) ,
\]

10
where \cite{21}

\[
|V_{ub}|_{\text{SM-exp}}^{\text{incl}} = (4.11 \pm 0.28) \times 10^{-3} .
\] (51)

The inclusive determination from $B \to \pi \ell \nu$, where only the vector current appears, leads to

\[
|V_{ub}|_{\text{SM-exp}}^{B \to \pi} = |V_{ub} + \epsilon_R \tilde{V}_{ub}| = (3.38 \pm 0.36) \times 10^{-3} ,
\] (52)

where the experimental value is taken from Ref. \cite{21}. Finally, a constraint on the $b \to u$ axial current can be obtained from the rare leptonic decay $B \to \tau \nu$. Using the theoretical expression

\[
\mathcal{B}(B \to \tau \nu)_{\text{SM}} = \frac{G_F^2 m_B m_{\tau}^2}{8\pi} \left(1 - \frac{m_{\tau}^2}{m_B^2}\right)^2 f_B^2 |V_{ub}^{\text{SM}}|^2 \tau_B ,
\] (53)

the experimental result $\mathcal{B}(B \to \tau \nu)^{\text{exp}} = (1.73 \pm 0.34) \times 10^{-4}$ \cite{25}, and $f_B = (192.8 \pm 9.9)$ MeV \cite{26}, we get

\[
|V_{ub}|_{\text{SM-exp}}^{B \to \tau} = |V_{ub} - \epsilon_R \tilde{V}_{ub}| = (5.14 \pm 0.57) \times 10^{-3} .
\] (54)

As noted first in \cite{12}, here the situation is very favourable for the contribution of RH currents, since the axial and vector exclusive determinations are substantially above and below the inclusive one (where the interference term is negligible). Performing a global fit to $V_{ub}$ and $\epsilon_R \tilde{V}_{ub}$ using the three constraints we get

\[
|V_{ub}| = (4.1 \pm 0.2) \times 10^{-3} , \quad \epsilon_R \text{ Re} \left(\frac{\tilde{V}_{ub}}{V_{ub}}\right) = -0.19 \pm 0.07 ,
\] (55)

with a correlation $\rho = -0.13$, namely an evidence of about 2.7$\sigma$ of a non-vanishing RH current contribution. In this case the quality of the fit is excellent ($\chi^2 \approx 0$) and substantially better than in the absence of RH currents. Most importantly the presence of right-handed currents removes the visible discrepancies between the various determinations, as shown in Fig. 1 (the SM case corresponds to the top of the vertical axis, where the three determinations of $|V_{ub}|$ are clearly different).

The above result has been obtained expanding to first order in $\epsilon_R$: in this limit we are sensitive only to the combination $\text{Re}(\tilde{V}_{ub}/V_{ub})$. Given the non-vanishing result for the RH term, we tried also a three-parameter fit, where also $\text{Im}(\tilde{V}_{ub}/V_{ub})$ is free to vary. In this case no unambiguous result is found, unless additional conditions are imposed. Imposing the condition $|\tilde{V}_{ub}| < |V_{ub}|$, the best solution is still the one in Eq. (55), which holds for a large interval of $\text{Im}(\tilde{V}_{ub}/V_{ub})$ around zero. Varying the phase of $\tilde{V}_{ub}/V_{ub}$ in a conservative range leads to

\[
|\epsilon_R \tilde{V}_{ub}| = (1.0 \pm 0.4) \times 10^{-3} , \quad \text{for} \quad -\frac{\pi}{4} < \arg \left(\frac{\tilde{V}_{ub}}{V_{ub}}\right) < \frac{\pi}{4} .
\] (56)

### 4.3 Global fit of the right-handed mixing matrix

The information we have collected in the previous section can be summarized as follows

\[
|\tilde{V}| \sim \begin{pmatrix}
< 1.4 & < 1.4 & 1.0 \pm 0.4 \\
- & - & < 2.0 \\
- & - & -
\end{pmatrix} \times \left(\frac{10^{-3}}{\epsilon_R}\right) ,
\] (57)
Figure 1: Constraints on $|V_{ub}|$ and $\epsilon_R \text{Re}(\tilde{V}_{ub}/V_{ub})$ from $d B \to \pi \ell \nu$ (green), $B \to X_u \ell \nu$ (blue), and $B \to \tau \nu$ (orange). The bands denote the ±1σ intervals of the various experimental constraints. The ellipse denotes the 1σ region of our best-fit solution.

where the inequalities correspond to the ±1σ interval and we have assumed small phases except for $\tilde{V}_{ub}$. The entries without figures have very weak direct experimental constraints. Altogether the constraints in Eq. (57) seem to be rather weak. However, thanks to the unitarity condition, they are sufficient to draw a series of interesting conclusions.

**Constraints on $\epsilon_R$.** The large value of $|\tilde{V}_{ub}|$ allows us to derive a significant constraint on the value of $\epsilon_R$ from the unitarity of the first row:

$$|\epsilon_R| = \left( |\epsilon_R \tilde{V}_{ad}|^2 + |\epsilon_R \tilde{V}_{us}|^2 + |\epsilon_R \tilde{V}_{ub}|^2 \right)^{1/2} = (1.0 \pm 0.5) \times 10^{-3}.$$ (58)

Given the bound on $\epsilon_L$ derived in Eq. (35), the possibility of $\epsilon_L$ and $\epsilon_R$ of the same order is perfectly allowed. Note also that the central value in Eq. (58) is in good agreement with the naive estimate of models with strong electroweak symmetry breaking, where we expect $c_{L,R} = O(1)$ and $\Lambda = 4\pi v \approx 3$ TeV.

We have no information to disentangle the sign of $\epsilon_R$ and $\tilde{V}_{ub}$. For simplicity in the following we assume $\epsilon_R$ to be positive. This assumption will not have any consequence
for our subsequent analysis since in all observables we always have a similar freedom in moving an overall sign from \( \tilde{V} \) to the Wilson coefficient of the effective operators.

**Constraints on** \( |\tilde{V}_{tb}|, |\tilde{V}_{ts}|, \text{ and } |\tilde{V}_{td}| \). Adopting the general parametrization in Eq. (12) we find a wide range for the three mixing angles compatible with data. The best-fit solution, obtained using only the constraints on \( \epsilon_R|\tilde{V}_{us}|, \epsilon_R|\tilde{V}_{ub}|, \epsilon_R|\tilde{V}_{ts}|, \epsilon_R|\tilde{V}_{td}| \), and \( \epsilon_R|\tilde{V}_{cb}| \), collected in the previous section, is

\[
\tilde{V} \sim \begin{pmatrix}
0 & -0.76 & -0.65 \\
0.88 & -0.31 & 0.36 \\
0.48 & 0.57 & -0.67
\end{pmatrix}.
\] (59)

However, other solutions with a rather different mixing structure are also compatible with data. For instance, the hierarchical scenario with \( \hat{s}_{13} \approx 1 \) and \( \hat{s}_{12} \approx \hat{s}_{23} \approx 0 \), provides also a good fit. What is interesting is that in all cases we can derive non-trivial constraints on the elements of the third row, that play a significant role in neutral-current observables (to be discussed in the next sections).

From the unitarity condition on the third column, and the large value of \( |\tilde{V}_{tb}| \), there follows a significant constraint on the maximal value of \( |\tilde{V}_{tb}| \). A large value of \( |\tilde{V}_{tb}| \) is particularly welcome since: i) it minimizes the values of \( |\tilde{V}_{ts}| \) and \( |\tilde{V}_{td}| \), that can induce too large contributions to FCNCs; ii) it maximizes the impact of right-handed currents in \( Z \rightarrow b\bar{b} \), which could help to improve the agreement with experiments (see Sect. 7.2).

In the following we thus concentrate on the scenario of maximal \( |\tilde{V}_{tb}| \). This is achieved with the ansatz

\[
\tilde{V}_0^{(I)} = \begin{pmatrix}
\hat{c}_{12}\hat{c}_{13} & \hat{s}_{12}\hat{c}_{13} & \hat{s}_{13} \\
-\hat{s}_{12} & \hat{c}_{12} & 0 \\
-\hat{c}_{12}\hat{s}_{13} & -\hat{s}_{12}\hat{s}_{13} & \hat{c}_{13}
\end{pmatrix},
\] (60)

which corresponds to the \( \hat{c}_{23} \rightarrow 1 \) limit of the general form in Eq. (12). Using the ansatz (60), and performing a global fit, we find that at the 90% C.L.

\[
|\tilde{V}_{tb}| < 0.73,
\] (61)

where the maximal value is obtained for \( \hat{s}_{13} \approx -0.68 \) and \( \hat{s}_{12} \approx -0.84 \). The fact that \( |\tilde{V}_{tb}| \) cannot reach 1 (as in the CKM) implies, in turn, that \( |\tilde{V}_{ts}| \) and \( |\tilde{V}_{td}| \) cannot be both vanishing. However, the sensitivity of the fit to \( \hat{s}_{12} \), which controls their relative strength, is very mild. The maximal \( |\tilde{V}_{tb}| \) in Eq. (61) is obtained for \( |\tilde{V}_{ts}| \approx |\tilde{V}_{td}| \approx 0.5 \), while if we require either \( |\tilde{V}_{ts}| \) or \( |\tilde{V}_{td}| \) to be vanishing, then the maximal allowed value of \( |\tilde{V}_{tb}| \) decreases by less than 10%. Only large positive values of \( \hat{s}_{12} \) are (slightly) disfavored by the \( \tilde{V}_{us} \) constraint in (37), if \( \hat{c}_{13} \) is positive.

In summary, we find that the scenario with maximal \( |\tilde{V}_{tb}| \) is well described by the ansatz (60) with

\[
|\hat{c}_{13}| \approx -\hat{s}_{13} \approx 0.7, \quad \epsilon_R \approx 1 \times 10^{-3},
\] (62)

and free \( \hat{s}_{12} \), provided that \( \text{sgn}(\hat{c}_{13}\hat{s}_{12}) = -1 \). In other words, we find a good description with

\[
\tilde{V}_0^{(II)} = \begin{pmatrix}
\pm\hat{c}_{12}\frac{\sqrt{2}}{2} & \pm\hat{s}_{12}\frac{\sqrt{2}}{2} & -\frac{\sqrt{2}}{2} \\
-\hat{s}_{12} & \hat{c}_{12} & 0 \\
\hat{c}_{12}\frac{\sqrt{2}}{2} & \hat{s}_{12}\frac{\sqrt{2}}{2} & \pm\frac{\sqrt{2}}{2}
\end{pmatrix}.
\] (63)
In the following we assume the matrix $\widetilde{V}_0^{\text{(II)}}$ as reference structure for the analysis of FCNCs. As we will discuss in Section 6, this structure is a key ingredient to generate a sizable non-standard contribution to $S_{\psi\phi}$. As a result, after we require a large $S_{\psi\phi}$ (as indicated by recent CDF [27] and D0 [28, 29] results), most of our conclusions will not depend on the choice of this ansatz. It should also be stressed that, contrary to the CKM case, having a zero in $\widetilde{V}_0^{\text{(II)}}$ does not prevent non-vanishing CP-violating effects thanks to the extra phases in Eq. (13).

As far as FCNCs are concerned, the only significant alternative to the ansatz $\widetilde{V}_0^{\text{(II)}}$ is the possibility of a vanishing small $|\widetilde{V}_{tb}|$, that we can achieve expanding around $\widetilde{c}_{13} \ll 1$ in (60):

$$\widetilde{V}_0^{\text{(III)}} = \begin{pmatrix} \widetilde{c}_{13}\widetilde{c}_{12} & \widetilde{c}_{13}\widetilde{s}_{12} & -1 \\ -\widetilde{s}_{12} & \widetilde{c}_{12} & 0 \\ \widetilde{c}_{12} & \widetilde{s}_{12} & \widetilde{c}_{13} \end{pmatrix}.$$  \hspace{0.5cm} (64)

As can be seen, this structure is very efficient in escaping all bounds from charged currents but for $b \rightarrow u$ transitions, where it has a maximal impact. However, it also implies naturally small effects in $B$ physics (both in $B_s$ and $B_d$ meson-anti-meson mixing, as well as on rare $B$ decays) and potentially large effects in $K$ physics (which are not allowed by data). We thus consider it less interesting with respect to the ansatz $\widetilde{V}_0^{\text{(II)}}$ in (63).

4.4 Summary

We conclude this section with a short summary of the results obtained from the analysis of charged currents:

- RH charged currents can help to reduce the tension between inclusive and exclusive determinations in $|V_{ub}|$, as recently pointed out in [12]. However, contrary to [12] and in agreement with [15], we find that RH charged currents do not have a significant impact in the determination of $|V_{cb}|$.

- The size of the RH charged-current operators necessary to solve the $|V_{ub}|$ problem points towards an effective new-physics scale $\Lambda \approx 3$ TeV.

- Thanks to unitarity, the full structure of the RH mixing matrix is quite constrained from a global fit of the available constraints. In particular, $|\widetilde{V}_{ub}|$ is constrained strongly by $|\widetilde{V}_{ub}|$ and unitarity, which in turn implies non-negligible contributions to FCNCs and meson-antimeson mixing through $|\widetilde{V}_{ts}|$ and $|\widetilde{V}_{td}|$.

- The two representative structures of the RH mixing matrix naturally emerge in view of the analysis of FCNCs: the structure $\widetilde{V}_0^{\text{(II)}}$ in (63) is particularly interesting, since it could allow large effects in the $B_s$ system, without being automatically excluded by the tight constraints from the $K$ and $B_d$ systems; the structure $\widetilde{V}_0^{\text{(III)}}$ in (64) would imply sizable effects only in the kaon system.

5 Dimension-six operators beyond charged currents

Having constrained the structure of the new mixing matrix from charged current processes, we are now ready to analyse the effects generated within our effective theory in neutral-current
processes. In particular in the next two sections we analyse the effects in the following set of theoretically clean observables: down-type particle-antiparticle mixing ($\varepsilon_K$ and $B_{d,s}$-$\bar{B}_{d,s}$ mixing), rare FCNC decays of $B$ and $K$ mesons with a lepton pair in the final state, and $Z \to b\bar{b}$.

The dimension-six effective operators contributing to these processes can be constructed by appropriate combinations of the bilinear structures in Eqs. (17)–(19). Rather than presenting lengthy expressions containing all possible operators, we limit ourself to analyse the impact of the most representative ones. In particular, we focus our attention on $\Delta F = 2$ operators built in terms of the $\bar{Q}R Y_u^\dagger Y_u\gamma^\mu Q_R$ bilinear, and $\Delta F = 1$ operators generating an effective right-handed flavour non-universal couplings of the $Z$ boson to down-type quarks.

5.1 $\Delta F = 2$ processes

The complete set of gauge-invariant dimension-six operators contributing to down-type $\Delta F = 2$ amplitudes, with the minimum number of Yukawa spurions, is

$$O^{(6)}_{LL} = [\bar{Q}_i^L (Y_u^\dagger Y_u^\dagger)_{ij} \gamma^\mu Q_j^L]^2,$$
$$O^{(6)}_{RR} = [\bar{Q}_i^R (Y_u^\dagger Y_u)_{ij} \gamma^\mu Q_j^R]^2,$$
$$O^{(6)}_{LR} = [\bar{Q}_i^L (Y_u^\dagger Y_u)_{ij} \gamma^\mu Q_j^L][\bar{Q}_i^R (Y_u^\dagger Y_u)_{ij} \gamma^\mu Q_j^R].$$

The first operator, which is present both in the general MFV framework [14] and in its constrained version [30] has been widely analysed in the literature. This operator generates short-distance corrections which have exactly the helicity structure and CKM factors of the SM short-distance terms. As a result, it cannot modify the SM predictions for the time-dependent CP asymmetries in $B_d$ and $B_s$ decays [30], in particular $S_{\psi K}$ and $S_{\psi \phi}$, respectively.

A much richer phenomenology is expected from $O^{(6)}_{RR}$ and $O^{(6)}_{LR}$, which are not present in the MFV framework and through which the new RH mixing matrix enters the game. For this reason, in the following we focus our attention only on these two operators, considering the following $\Delta F = 2$ effective Lagrangian:

$$\mathcal{L}^{\Delta F=2} = \frac{c_{RR}}{\Lambda^2} O^{(6)}_{RR} + \frac{c_{LR}}{\Lambda^2} O^{(6)}_{LR}.$$

Here $c_{RR}$ and $c_{LR}$ are flavour-blind dimension-less coefficients, whose size will be discussed below.

5.2 $\Delta F = 1$ processes

The list of operators relevant to $\Delta F = 1$ FCNC processes with a lepton pair in the final state can be divided into three categories: 1) operators with two quarks and two Higgs fields; 2) operators with two quarks and two lepton fields; 3) dipole-type operators with two quarks and one SM gauge field.

In the first class we have the following two operators

$$O^{(6)}_{R_{z21}} = i \bar{Q}_R^i (Y_u^\dagger Y_u)_{ij} \gamma^\mu H^\dagger D_\mu HQ_j^L,$$
$$O^{(6)}_{R_{z22}} = i \bar{Q}_R^i (Y_u^\dagger Y_u)_{ij} \gamma^\mu \tau_i Q_j^R Tr \left( H^\dagger D_\mu H \tau^i \right).$$

15
and the corresponding LH operators obtained from (69) through \( Q_R \rightarrow Q_L \) and \( Y_u \leftrightarrow Y_u^\dagger \). The latter are of MFV type and have already been analysed in the literature. On the other hand, the two operators in (69) give rise to an effective RH coupling of the \( Z \) boson of the type \( \bar{d}_R^i \gamma^\mu d_R^j Z_\mu \) which is not present in the MFV framework. This coupling is particularly interesting since it allows us to establish connections between RH effects in \( Z \rightarrow b\bar{b} \) and in rare \( K \) and \( B \) decays. In the following we will analyse these connections by means of the effective Lagrangian

\[
L_{\Delta F=1} = \frac{c_{RZ1}}{\Lambda^2} O^{(6)}_{RZ1} + \frac{c_{RZ2}}{\Lambda^2} O^{(6)}_{RZ2}.
\]

As far as the other two classes of \( \Delta F = 1 \) operators are concerned: the operators with two quarks and two lepton fields do not lead to effects qualitatively different that those obtained after integrating out the \( Z \) in Eq. (69). The left-right operators with the photon field are tightly constrained by \( B \to X_s \gamma \) and, similarly to the MFV case, once the \( B \to X_s \gamma \) constraint is imposed they do not lead to significant effects in other processes.

6 Meson anti-meson mixing

6.1 Preliminaries

The presence of the right-handed currents in addition to the left-handed ones present in the SM can have considerable impact on particle-antiparticle mixing dominantly through the generation of the LR operators that renormalize strongly under QCD and in the case of \( K^0 - \bar{K}^0 \) mixing have chirally enhanced hadronic matrix elements. Below we present the general structure of the mixing amplitude \( M_{12} \) from which the observables like \( \varepsilon_K \) and the CP-asymmetries \( S_{\psi K} \) and \( S_{\psi \phi} \) can be derived. Similarly \( \Delta M_d \) and \( \Delta M_s \) can be calculated.

6.2 Effective Hamiltonian for \( \Delta S = 2 \) transitions

We will illustrate the full procedure on the example of \( \Delta S = 2 \) transitions. The complete set of operators of dimensions six involved in the presence of left-handed and right-handed currents consists of the following operators [31]:

\[
\begin{align*}
Q_{1VLL} &= (\bar{s}_L \gamma_\mu d_L) (\bar{s}_L \gamma^\mu d_L), \\
Q_{1VRR} &= (\bar{s}_R \gamma_\mu d_R) (\bar{s}_R \gamma^\mu d_R), \\
Q_{1LR} &= (\bar{s}_L \gamma_\mu d_L) (\bar{s}_R \gamma^\mu d_R), \\
Q_{2LR} &= (\bar{s}_R d_L) (\bar{s}_L d_R),
\end{align*}
\]

where we suppressed colour indices as they are summed up in each factor \( P_{L,R} \).

In the SM only \( Q_{1VLL} \) contributes. However, in the presence of right-handed currents three additional operators have to be considered. \( Q_{1VRR} \) renormalizes under QCD, similarly to \( Q_{1VLL} \), fully independently of other operators and in fact the renormalization group running for the Wilson coefficient of \( Q_{1VRR} \) is identical to the one of \( Q_{1VLL} \); QCD does not care about the sign of \( \gamma_5 \). On the other hand \( Q_{1LR} \) and \( Q_{2LR} \) mix under renormalization.

Let us next assume that at some scale \( \mathcal{O}(1 \text{ TeV}) \), to be denoted by \( \mu_R \), the new physics (NP) is integrated out. In the absence of QCD corrections the effective Hamiltonian for
\[ \Delta S = 2 \text{ corresponding to } \mu_R = \mathcal{O}(\Lambda) \text{ and including only NP contributions reads} \]
\[
\left[ H_{\text{eff}}^{\Delta S=2}(\mu_R) \right]_{\text{NP}} = \frac{1}{\Lambda^2} \left[ C_1^{VLL}(\mu_R, K) Q_1^{VLL} + C_1^{VRR}(\mu_R, K) Q_1^{VRR} + C_1^{LR}(\mu_R, K) Q_1^{LR} \right].
\] (72)

The following comments should be made:

- The coefficients \( C_a^i \) with \( i = 1, 2 \) and \( a = VLL, VRR, LR \) depend generally on the system considered \((K, B_{s,d})\) as is the case also of the SM.
- At the scale \( \mu_R \), before integrating out the \( W_L^+ \) fields, the Wilson coefficient \( C_1^{VLL} \) receives only contributions from the NP. Running down to low scales, \( C_1^{VLL} \) encodes both SM and NP contributions, as discussed below.
- The coefficient \( C_2^{LR}(\mu_R, K) \) vanishes at this high scale in the absence of QCD effects and it is also vanishing at this scale in the LO renormalization group (RG) analysis. However, at NLO it is generally \( \mathcal{O}(\alpha_s(\mu_R)) \).

Next usually (72) is evolved by RG down to low energy scales. In this process \( C_2^{LR} \) becomes non-vanishing even in the LO approximation. As \( Q_1^{VRR} \) and the complex \((Q_1^{LR}, Q_2^{LR})\) do not mix with \( Q_1^{VLL} \) and the RG evolution of the latter operator can be split into SM and NP part, the NP contribution to the low energy effective Hamiltonian at a low scale \( \mu_K \) can be immediately evaluated by means of analytic expressions in [31].

While working with Wilson coefficients and operator matrix elements at low energy scales is a common procedure, it turns out that for phenomenology it is more useful to work directly with \( C_i(\mu_R, K) \) and with the hadronic matrix elements of the corresponding operators also evaluated at this high scale. The latter matrix elements are given by [31]
\[
\langle \bar{K}_0 | Q_i^a | K_0 \rangle = \frac{2}{3} M_{KK}^2 F_K^2 P_i^a(K),
\] (73)

where the coefficients \( P_i^a(K) \) collect compactly all RG effects from scales below \( \mu_R \) as well as hadronic matrix elements obtained by lattice methods at low energy scales.

With this information at hand we can present compactly basic formulae for \( \Delta S = 2 \) and \( \Delta B = 2 \) observables of interest.

### 6.3 Basic formulae for \( \Delta S = 2 \) observables

The off-diagonal element in \( K^0 - \bar{K}^0 \) mixing \( M_{12}^K \) can then be decomposed into SM and NP parts (RH current contributions in our case)
\[
M_{12}^K = (M_{12}^K)_{\text{SM}} + (M_{12}^K)_{\text{NP}},
\] (74)

where the SM contribution can be found, for instance, in Ref. [31], and the NP part is obtained from
\[
2m_K (M_{12}^K)_{\text{NP}}^* = \langle \bar{K}^0 | [H_{\text{eff}}^{\Delta S=2}(\mu_R)]_{\text{NP}} | K^0 \rangle.
\] (75)

Using (72), (73) and (75) we find
\[
(M_{12}^K)_{\text{NP}} = \frac{1}{3\Lambda^2} m_K F_K^2 \cdot \left[ (C_1^{VLL}(\mu_R, K) + C_1^{VRR}(\mu_R, K)) P_1^{VLL}(K) + C_1^{LR}(\mu_R, K) P_1^{LR}(K) \right]^*,
\] (76)
with
\[ P_{1}^{VLL}(K) \approx 0.50, \quad P_{1}^{LR}(K) \approx -52 \] (77)
obtained by means of analytic formulae in [31], with the hadronic matrix elements from [32], the updated numerical inputs in Table 1, and the matching scale \( \mu_R = 1.5 \text{ TeV} \).

The \( KL - KS \) mass difference and the CP-violating parameter \( \varepsilon_K \) are given respectively
\[
\Delta M_K &= 2 \Re M_{12}^K + (\Delta M_K)_{LD}, \\
\varepsilon_K &= \frac{\kappa_e e^{i\varphi_e}}{\sqrt{2} (\Delta M_K)_{exp}} \Im M_{12}^K,
\] (78)
where \( \varphi_e = (43.51 \pm 0.05)^\circ \) takes into account that \( \varphi_e \neq \pi/4 \) and \( \kappa_e = 0.94 \pm 0.02 \) includes an additional effect from long-distance (LD) contributions (for a recent detailed analysis of \( \varepsilon_K \) within the SM see [33, 34]).

Matching the RH effective Lagrangian defined in [68] to the general effective \( \Delta S = 2 \) Hamiltonian in (72), the Wilson coefficients at the high scale read
\[
C_1^{VLL}(\mu_R, K) = 0,
\] (79)
where the terms on the right-hand side are obtained employing the structure (63) for the RH mixing matrix.

The non-standard contributions to \( \Delta S = 2 \) amplitudes are exceedingly large compared to the SM term (and compared with data) unless the Wilson coefficients \( c_{RR} \) and \( c_{LR} \) or one of the two mixing terms \( \tilde{c}_{12} \) or \( \tilde{s}_{12} \) are very small. By construction, \( c_{RR} \) and \( c_{LR} \) are flavour-blind and therefore the same in the \( B_d \) and \( B_s \) system. On the other hand, the \( \tilde{c}_{12} \) and \( \tilde{s}_{12} \) dependencies in the three systems considered are non-universal, as seen in Table 2 with the observables in the \( K \)-mixing, \( B_d \) mixing and \( B_s \)-mixing dominated by \( \tilde{c}_{12}\tilde{s}_{12} \), \( \tilde{c}_{12} \) and \( \tilde{s}_{12} \), respectively. Since both \( \Delta S = 2 \) and \( B_d \) mixing are strongly constrained, and the

| parameter | value     | parameter | value     |
|-----------|-----------|-----------|-----------|
| \( F_K \) | \( (155.8 \pm 1.7) \text{ MeV} \) | \( \Delta M_K \) | \( (5.292 \pm 0.009) \times 10^{-3} \text{ ps}^{-1} \) |
| \( F_{B_d} \) | \( (192.8 \pm 9.9) \text{ MeV} \) | \( \Delta M_d \) | \( (0.507 \pm 0.005) \text{ ps}^{-1} \) |
| \( F_{B_s} \) | \( (238.8 \pm 9.5) \text{ MeV} \) | \( \Delta M_s \) | \( (17.77 \pm 0.12) \text{ ps}^{-1} \) |
| \( B_K \) | \( 0.725 \pm 0.026 \) | \( \vert V_{tb} \vert \) | \( 1 \pm 0.06 \) |
| \( B_{B_d} \) | \( 1.26 \pm 0.11 \) | \( \vert V_{td} \vert \) | \( (8.3 \pm 0.5) \times 10^{-3} \) |
| \( B_{B_s} \) | \( 1.33 \pm 0.06 \) | \( \vert V_{ts} \vert \) | \( 0.040 \pm 0.003 \) |
| \( M_K \) | \( 0.497614 \text{ GeV} \) | \( \sin(2\beta_s) \) | \( 0.038 \pm 0.003 \) |
| \( M_{B_d} \) | \( 5.2795 \text{ GeV} \) | \( \gamma \) | \( 1.09 \pm 0.12 \) |
| \( M_{B_s} \) | \( 5.3664 \text{ GeV} \) | \( \varepsilon_{K}^{exp} \) | \( (2.29 \pm 0.01) \times 10^{-3} \) |
| \( m_t(m_t) \) | \( (163.5 \pm 1.7) \text{ GeV} \) | \( \varepsilon_{K}^{exp} \) | \( S_{\psi K}^{exp} \) | \( 0.672 \pm 0.023 \) |

Table 1: Values of the input parameters used in our analysis of \( \Delta F = 2 \) processes.
For the RH matrix are those in Eq. (63) and (64), respectively.

factors expanded in powers of $\lambda$ the SM (LH sector) and in the RH sector. In the SM case approximate expressions of the CKM

Table 2: Mixing structures relevant to the three down-type $\Delta F = 2$ and FCNC amplitudes in the SM (LH sector) and in the RH sector. In the SM case approximate expressions of the CKM factors expanded in powers of $\lambda = |V_{us}|$ are also shown. In the RH case the two parametrizations for the RH matrix are those in Eq. (63) and (64), respectively.

| Mixing term | Matrix | $s \rightarrow d$ | $b \rightarrow d$ | $b \rightarrow s$ |
|-------------|--------|------------------|------------------|------------------|
| $V_{ti}^*V_{tj}$ | CKM | $V_{ti}^*V_{td} \approx -\lambda^5 e^{-i\beta}$ | $V_{tb}^*V_{td} \approx \lambda^3 e^{-i\beta}$ | $V_{ts}^*V_{ts} \approx -\lambda^2 e^{-i\beta}$ |
| $\tilde{V}_{ti}^*\tilde{V}_{tj}$ | $\tilde{V}_0^{(II)}$ | $\frac{1}{2} \tilde{c}_{12} \tilde{s}_{12} e^{i(\phi_d^d - \phi_d^s)}$ | $\pm \frac{1}{2} \tilde{c}_{12} \tilde{c}_{12} e^{i(\phi_d^d - \phi_d^s)}$ | $\pm \frac{1}{2} \tilde{s}_{12} e^{i(\phi_d^d - \phi_d^s)}$ |
| | $\tilde{V}_0^{(III)}$ | $\tilde{c}_{12} \tilde{s}_{12} e^{i(\phi_d^d - \phi_d^s)}$ | $\tilde{c}_{12} \tilde{c}_{13} e^{i(\phi_d^d - \phi_d^s)}$ | $\tilde{s}_{12} \tilde{c}_{13} e^{i(\phi_d^d - \phi_d^s)}$ |

data from CDF and D0 give some hints for sizable NP contributions in the $B_s$ mixing, it is natural to assume in both scenarios for $\tilde{V}$ that $\tilde{c}_{12} \ll 1$. In this limit the non-vanishing Wilson coefficients relevant for $K^0 - \bar{K}^0$ mixing and $\varepsilon_K$ at the high scale read with (63)

\[
C_1^{V_{RR}}(\mu_R, K) \approx -\frac{c_{RR}}{4} y_t^4 e^{2i\phi_d^d} \tilde{c}_{12}^2, \\
C_1^{LR}(\mu_R, K) \approx -\frac{c_{LR}}{2} y_t^4 e^{i\phi_d^d} V_{ts}^*V_{td} \tilde{c}_{12},
\]

where $\phi_{21}^d = (\phi_d^d - \phi_d^s)$. Using these expressions in the above formulae for $\Delta M_K$ and $\varepsilon_K$, and taking into account the numerical inputs in Table 1 we obtain:

\[
(\Delta M_K)_{RH} = (\Delta M_K)_{exp} \times \left[ -2.5 \times 10^4 \times c_{RR} \tilde{c}_{12}^2 \cos(2\phi_{21}^d) \\
- 1.7 \times 10^3 \times c_{LR} \tilde{c}_{12} \cos(\phi_{21}^d - \beta - \beta_s) \right] \frac{3 \text{ TeV}^2}{\Lambda^2}, \tag{81}
\]

\[
(\varepsilon_K)_{RH} = |\varepsilon_K|_{exp} e^{i\phi_0} \times \left[ 3.7 \times 10^6 \times c_{RR} \tilde{c}_{12}^2 \sin(2\phi_{21}^d) \\
+ 2.5 \times 10^5 \times c_{LR} \tilde{c}_{12} \sin(\phi_{21}^d - \beta - \beta_s) \right] \frac{3 \text{ TeV}^2}{\Lambda^2}, \tag{82}
\]

where, as usual, $\beta$ and $\beta_s$ denote the phases of $V_{td}$ and $V_{ts}$ in the standard CKM convention:

\[
V_{td} = |V_{td}| e^{-i\beta} \quad \text{and} \quad V_{ts} = -|V_{ts}| e^{-i\beta}. \tag{83}
\]

As can be seen, the RH contribution is potentially very large and, independently of the possible value of the CP-violating phase $\phi_{21}^d$, we get strong constraints on the combinations $c_{RR} \tilde{c}_{12}^2$ and $c_{LR} \tilde{c}_{12}$. In particular, even assuming a vanishing $\phi_{21}^d$, we get

\[
c_{RR} \tilde{c}_{12}^2 < 2.0 \times 10^{-5}, \quad \text{from} \quad (\Delta M_K)_{RH} < 0.5(\Delta M_K)_{exp}, \tag{84}
\]

\[
c_{LR} \tilde{c}_{12} < 1.0 \times 10^{-6}, \quad \text{from} \quad |\varepsilon_K|_{RH} < 0.1|\varepsilon_K|_{SM}. \tag{85}
\]

While the bound on $c_{RR} \tilde{c}_{12}^2$ can only become stronger for non-vanishing values of $\phi_{21}^d$, the constraint on $c_{LR} \tilde{c}_{12}$ can be relaxed to $3 \times 10^{-4}$ in the fine-tuned scenario where $\phi_{21}^d$ cancels exactly the CKM phase of $V_{ts}^*V_{td}$. As we will show in the next section, these constraints imply negligible contribution of RH currents to $B_d$ mixing, which is one of the important results of our paper.
Before analysing $\Delta B = 2$ observables, we briefly discuss what happens if we do not employ the ansatz (63) for the RH mixing matrix and, in particular, if we do not assume $\tilde{c}_{12} \ll 1$. As shown in Table 2 employing the structure (64) for the RH matrix the mixing structures relevant to the kaon system change only by a factor of two. As a result, the corresponding bounds on the mixing terms are obtained from (84)–(85) with the replacement $\tilde{c}_{12} \rightarrow 2\tilde{c}_{12}\tilde{s}_{12}$.

In this case we can escape the kaon bounds and have sizable effects in $B_d$ mixing if $\tilde{c}_{13}$ is not too small. However, the key ingredient for sizable effects in $B_d$ mixing is $\tilde{s}_{12} \ll 1$, a configuration that necessarily imply small effects in $B_s$ mixing.

### 6.4 Basic formulae for $\Delta B = 2$ observables

Similarly to the kaon system, for $B_{d,s}^0 - \bar{B}_{d,s}^0$ mixing we can write

$$M_{12}^q = (M_{12}^q)_{\text{SM}} + (M_{12}^q)_{\text{NP}}$$

where $q = d, s$. For the NP contribution we find

$$(M_{12}^q)_{\text{NP}} = \frac{1}{3\Lambda^2} m_{B_q} F_{B_q}^2 \left[ (C_1^{VLL}(\mu_R, B) + C_1^{VRR}(\mu_R, B)) P_1^{VLL}(B) \right. + \left. C_1^{LR}(\mu_R, B) P_1^{LR}(B) \right]^*,$$

with

$$P_1^{VLL}(B) \approx 0.70, \quad P_1^{LR}(B) \approx -3.2$$

obtained by means of analytic formulae in [31], with the hadronic matrix elements from [32], the updated numerical inputs in Table 1 and the matching scale $\mu_R = 1.5$ TeV.

For the mass differences in the $B_{d,s}^0 - \bar{B}_{d,s}^0$ systems we have

$$\Delta M_q = 2|M_{12}^q|,$$

$$M_{12}^q = (M_{12}^q)_{\text{SM}} C_{B_q} e^{2i\varphi_{B_q}}, \quad (q = d, s)$$

where

$$\begin{align*}
(M_{12}^d)_{\text{SM}} &= \left| (M_{12}^d)_{\text{SM}} \right| e^{2i\beta}, \\
(M_{12}^s)_{\text{SM}} &= \left| (M_{12}^s)_{\text{SM}} \right| e^{2i\beta_s},
\end{align*}$$

and $C_{B_q} \neq 1$ and $\varphi_{B_q} \neq 0$ summarize the NP effects.

We find then

$$\Delta M_q = (\Delta M_q)_{\text{SM}} C_{B_q}$$

and

$$S_{\psi K_S} = \sin(2\beta + 2\varphi_{B_q}), \quad S_{\psi \phi} = \sin(2|\beta_s| - 2\varphi_{B_s}),$$

with the latter two observables being the coefficients of $\sin(\Delta M_{d,t})$ and $\sin(\Delta M_{s,t})$ in the time dependent asymmetries in $B_d^0 \rightarrow \psi K_S$ and $B_s^0 \rightarrow \psi \phi$, respectively. Thus in the presence of non-vanishing $\varphi_{B_d}$ and $\varphi_{B_s}$ these two asymmetries do not measure $\beta$ and $\beta_s$ but $(\beta + \varphi_{B_d})$ and $(|\beta_s| - \varphi_{B_s})$, respectively.
The non-vanishing Wilson coefficients at the high scale for the $B_{s,d}$ systems are

$$C_1^{VRR}(\mu_R, B_q) = -c_{RR} y_t^4 e^{2i(\phi_d^d - \phi_d^s)} \left[ (\bar{V}_0)_{ib}^* (\bar{V}_0)_{tq} \right]^2,$$

$$C_1^{LR}(\mu_R, B_q) = -c_{LR} y_t^4 e^{i(\phi_d^d - \phi_d^s)} V_{ib} V_{tq} (\bar{V}_0)_{ib}^* (\bar{V}_0)_{tq} .$$

(94)

Working in the limit $\tilde{c}_{12} \ll 1$ (hence $\tilde{s}_{12} \approx 1$) we get

$$C_1^{VRR}(\mu_R, B_d) \approx -\frac{c_{RR}}{4} y_t^4 e^{2i(\phi_d^d - \phi_d^s)} \tilde{c}_{12}^2,$$

$$C_1^{LR}(\mu_R, B_d) \approx \mp \frac{c_{LR}}{2} y_t^4 e^{i(\phi_d^d - \phi_d^s)} V_{ib} V_{td} \tilde{c}_{12},$$

$$C_1^{VLL}(\mu_R, B_{s,d}) = 0 ,$$

(95)

where the $\mp$ sign reflects the $\pm$ in [63]. Using these expressions in the formulae for $M_{12}^d$ we obtain

$$(M_{12}^d)_{SM+RH} = (M_{12}^d)_{SM} \times \left[ 1 + \left( -6.1 \times 10^3 \times c_{RR} \tilde{c}_{12}^2 e^{-2i(\phi_d^d + \beta)} \pm 4.7 \times 10^2 \times c_{LR} \tilde{c}_{12} e^{-i(\phi_d^d + \beta)} \right) \frac{(3 \text{ TeV})^2}{\Lambda^2} \right] ,$$

(96)

$$(M_{12}^s)_{SM+RH} = (M_{12}^s)_{SM} \times \left[ 1 + \left( -2.5 \times 10^2 \times c_{RR} e^{-2i(\phi_d^d + \beta_s)} \mp 0.9 \times 10^2 \times c_{LR} e^{-i(\phi_d^d + \beta_s)} \right) \frac{(3 \text{ TeV})^2}{\Lambda^2} \right] ,$$

(97)

where, similarly to $\phi_3^d$, we have defined $\phi_{3i}^d = (\phi_3^d - \phi_1^d)$. In obtaining the numerical values we have evaluated $(M_{12}^d)_{SM}$ using the inputs in Table 1.

As anticipated, given the bounds on $c_{RR} \tilde{c}_{12}^2$ and $c_{LR} \tilde{c}_{12}$ in [84]-[85], following from the neutral kaon system, in this framework non-standard contributions to both modulo and phase of $B_d$ mixing are safely negligible. On the other hand, sizable contributions to the $B_s$ system are possible if $c_{RR,LR}$ are in the $10^{-3}$--$10^{-2}$ range and $\tilde{c}_{12}$ is small enough to satisfy the kaon bounds.

If we do not assume $\tilde{c}_{12} \ll 1$ and, more generally, go beyond the ansatz [63] for the RH mixing matrix, sizable contributions to $B_d$ mixing are possible assuming $\tilde{s}_{12} \approx 1$. However, as already stated, this precludes the possibility of large effects in the $B_s$ system.

### 6.5 Combined fit of $\varepsilon_K$ and $B_s$ mixing

Here we discuss in more detail the interesting scenario with $c_{RR,LR}$ in the $10^{-3}$--$10^{-2}$ range and $c_{12} \ll 1$, where we can accommodate a large CP-violating phase in $B_s$ mixing, as hinted by CDF [27] and D0 [28,29] and, at the same time, satisfy the bounds from $\varepsilon_K$. Values of $c_{RR,LR}$ in the $10^{-3}$--$10^{-2}$ range are substantially lower than the $O(1)$ Wilson coefficients determined from charged-current interactions (assuming $\Lambda = 3$ TeV as reference scale). However, it is perfectly conceivable that the $\Delta F = 2$ operators are loop-suppressed with respect to the charged-current ones, such that a $10^{-3}$--$10^{-2}$ relative suppression of the corresponding Wilson coefficients can naturally be accommodated.

We stress that a large CP-violating phase in $B_s$ mixing is not a clear prediction of the model we are considering. As pointed out in the previous section, the only clear prediction...
is the absence of significant contributions to $B_d$ mixing, implied by the bounds from the kaon system, if we require a large CP-violating phase in $B_s$ mixing. Still, it is interesting to check if this experimental “anomaly” can be solved with reasonable values of the free parameters of the effective theory we are considering. To better investigate this point, we analyse separately the cases where the leading correction to the SM in $B_s$ mixing is induced by $c_{RR}$ and $c_{LR}$, respectively.

Assuming $c_{LR} \ll c_{RR}$ and neglecting the tiny contribution of $\beta_s$, the constraints of the $B_s$ system imply

$$\frac{(\Delta M_s)^{\text{SM+RH}}}{(\Delta M_s)^{\text{SM}}} = \left| 1 - 2.6 \times 10^2 \times c_{RR} e^{-2i\phi_2^d} \right| = \frac{(\Delta M_s)^{\text{exp}}}{(\Delta M_s)^{\text{SM}}} \approx 0.96 \pm 0.15 \ , \quad (98)$$

$$S_{\psi \phi} = - \frac{2.6 \times 10^2 \times c_{RR} \sin(2\phi_{32}^d)}{1 - 2.6 \times 10^2 \times c_{RR} e^{-2i\phi_2^d}} \approx 0.6 \pm 0.3 \ , \quad (99)$$

where the latter numerical entry is only a rough indicative value for the CP asymmetry favored by the Tevatron experiments (for more details see [27–29] and the model-independent analysis in [39]). The central values of these equations are full-filled with the following four-fold solution:

$$c_{RR} \approx \pm 7.3 \times 10^{-3} \quad \text{and} \quad \sin(2\phi_{32}^d) \approx \mp 0.30 \ ,$$

$$c_{RR} \approx \pm 2.3 \times 10^{-3} \quad \text{and} \quad \sin(2\phi_{32}^d) \approx \mp 0.95 \ . \quad (100)$$

As anticipated, these values are in good agreement with the naive expectation of a $1/(16\pi^2) \approx 6 \times 10^{-3}$ suppression between this $\Delta F = 2$ operator (presumably generated at the loop level) relative to those contributing to right-handed charged currents (presumably generated at the tree level).

Having fixed $c_{RR}$, we can ask which are the values of the mixing angles necessary to satisfy the bounds from the kaon system. To this end, it is first important to note that, due to the higher value of $|V_{ub}|$ determined from charged-current processes with the inclusion of RH currents, the prediction of $\varepsilon_K$ in this framework without extra contributions is in excellent agreement with data (contrary to what happens in the SM [33,40]). Indeed the value of $\sin(2\beta)$ determined by tree-level observables only, namely $|V_{ub}|$ and $\gamma$ (following the analysis in [37]) is

$$\sin(2\beta)_{\text{tree}}^{\text{RH}} = 0.77 \pm 0.05 \ . \quad (101)$$

This value is substantially higher than the corresponding result obtained in the SM, $\sin(2\beta)_{\text{tree}}^{\text{SM}} = 0.734 \pm 0.034$ [37], where the inclusive and exclusive determinations of $|V_{ub}|$ are averaged. As a result of this higher value of $\sin(2\beta)$, the tension between the experimental value of $\varepsilon_K$ and its prediction within the SM (see e.g. [35,41] for a recent analysis) is automatically solved.

Despite there is no need for non-standard contributions to $\varepsilon_K$, the theoretical errors on this observable allow for extra contributions within $\approx \pm 10\%$ of the SM amplitude. This condition is obtained for

$$|\tilde{c}_{12}| \sin(2\phi_{21}^d)|^{1/2} < 1.9 \times 10^{-3} \ , \quad \text{for} \ |c_{RR}| \approx 7.3 \times 10^{-3} ,$$

$$|\tilde{c}_{12}| \sin(2\phi_{21}^d)|^{1/2} < 3.4 \times 10^{-3} \ , \quad \text{for} \ |c_{RR}| \approx 2.3 \times 10^{-3} . \quad (102)$$

These values are small but not highly fine-tuned: for CP-violating phases of $O(0.1)$, the mixing angle $\tilde{c}_{12}$ can reach values of $O(10^{-2})$, which are larger than the CKM element $|V_{ub}|$.  

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Due to the large chiral enhancement of the contribution of $Q^{LR}_1$ to $\varepsilon_K$, more fine-tuning is required if $c_{LR}$ provides the dominant contribution to $B_s$ mixing. Indeed repeating the above argument for $c_{RR} \ll c_{LR}$ leads to the following four-fold solution from $B_s$ mixing,

$$c_{LR} \approx \pm 2.0 \times 10^{-2} \quad \text{and} \quad \sin(2\phi_{32}^d) \approx \mp 0.30,$$

$$c_{LR} \approx \pm 0.6 \times 10^{-2} \quad \text{and} \quad \sin(2\phi_{32}^d) \approx \mp 0.95,$$

and the following conditions from $\varepsilon_K$:

$$|\tilde{c}_{12} \sin(\phi_{21}^d - \beta + \beta_s)| < 0.2 \times 10^{-4}, \quad \text{for} \quad |c_{LR}| \approx 2.0 \times 10^{-2},$$

$$|\tilde{c}_{12} \sin(\phi_{21}^d - \beta + \beta_s)| < 0.6 \times 10^{-4}, \quad \text{for} \quad |c_{LR}| \approx 0.6 \times 10^{-2}.$$  

It is clear that in this case the condition on the 1-2 mixing is more stringent than in (102).

### 6.6 Summary

The main results of the $\Delta F = 2$ analysis can be summarized as follows:

- A large CP-violating phase in $B_s$ mixing, as hinted by the Tevatron experiments, can be accommodated for natural values of the free parameters. The two necessary ingredients for this mechanism to work are: 1) Wilson coefficients of $O(1/(16\pi^2))$ for the $\Delta F = 2$ operators in (68), assuming as reference scale $\Lambda = 3$ TeV; 2) a RH mixing matrix with the following from $\tilde{\mathbf{V}}(B_s$ mixing)$|^0 \approx$ 

$$\begin{pmatrix}
0 & \sqrt{2} & \sqrt{2} \\
1 & 0 & 0 \\
0 & \sqrt{2} & \sqrt{2}
\end{pmatrix},$$  

where the null entries should not be taken as exact zeros, but rather as very small entries.

- According to the RH mixing structure in (105), with improved experimental precision it should be possible to resolve the presence of RH currents in $s \to u$ charged-current transitions, or a $O(10^{-3})$ deviation in the determination of $|V_{us}|$ from $K \to \pi\ell\nu$ and $K \to \ell\nu$ decays.

- Thanks to the large value of $\sin(2\beta)$, following from the inclusion of RH currents in the determination on $|V_{ub}|$, the prediction of $\varepsilon_K$ in this framework is in excellent agreement with data without extra contributions.

- The combination of a large CP-violating phase in $B_s$ mixing and only small NP effects allowed by $\varepsilon_K$ implies negligible effects in $B_d$ mixing. This, in turn, implies a tension between the measured valued of $S_{\psi K}$ and the predicted value of $\sin(2\beta)$ in this framework. The only possibility to solve this problem is to assume $s_{12} \ll 1$, giving up the possibility of NP effects in $B_s$ mixing.

### 7 $Z$-mediated FCNCs and $Z \to b\bar{b}$

#### 7.1 Modification of the RH couplings of the $Z$ boson

As discussed in Section 5 in the $\Delta F = 1$ sector we focus our attention on the effective Lagrangian (70). The effective operators appearing in this Lagrangian are equivalent to
$O_{R_{a1}}^{(6)}$ and $O_{R_{b2}}^{(6)}$ analysed in Section 3.2 but for the additional insertion of the combination of Yukawa matrices $(Y_u^0 Y_u)_{ij}$. After the breaking of the electroweak symmetry, they lead to the following effective right-handed couplings of the $Z$ boson to down-type quarks:

$$L_{\text{eff}}^{(Z_R)} = -\frac{g}{c_W} v^2 \frac{(c_{R_{Z1}} + 2c_{R_{Z2}}) y_t^2 \tilde{V}_t \tilde{V}_{ij} d_R^\gamma d_R^\gamma} {2 \Lambda^2} Z_\mu ,$$

where $c_W = \cos \Theta_W$ (similarly, in the following we use $s_W = \sin \Theta_W$).

Denoting the effective couplings of the $Z$ to down-type quarks as follows

$$L_{\text{eff}}^Z = \frac{g}{c_W} \left( g_L^{ij} d_L^\gamma d_L^\gamma + g_R^{ij} d_R^\gamma d_R^\gamma \right) Z_\mu ,$$

the SM contribution, evaluated at the one-loop level in the 't Hooft-Feynman gauge in the large top-mass limit, is

$$(g_L^{ij})_{\text{SM}} = \left( -\frac{1}{2} + \frac{1}{3} s_W^2 \right) \delta_{ij} + \frac{g^2}{8 \pi^2} V_{ii}^* V_{jj} C_0(x_t) , \quad x_t = \frac{m_t^2}{m_W^2} , \quad (g_R^{ij})_{\text{SM}} = \frac{1}{3} s_W^2 \delta_{ij} .$$

The loop function $C_0(x_t)$, that in the large $x_t$ limit is gauge independent $(g^2 C_0(x_t) \rightarrow g^2 x_t/8 = g^2/4$ for $m_t^2 \gg m_W^2$), can be found in [42]. Using these notations, the effect of the RH operators $O_{R_{a1}}^{(6)}$ and $O_{R_{b2}}^{(6)}$ can be included as a modification of the RH effective coupling:

$$(g_R^{ij})_{\text{tot}} = (g_R^{ij})_{\text{SM}} + (\Delta g_R^{ij})_{RH} , \quad (\Delta g_R^{ij})_{RH} = -\frac{v^2 (c_{R_{Z1}} + 2c_{R_{Z2}}) y_t^2 \tilde{V}_t \tilde{V}_{ij}} {2 \Lambda^2}.$$ 

### 7.2 $Z \rightarrow b\bar{b}$

The experimental determination of the effective couplings of the $Z$ bosons to $b$ quarks resulting from the global fit of electroweak data collected by the LEP and the SLD experiments is [43]

$$(g_{b\bar{b}}^{(6)})_{\text{exp}} = -0.4182 \pm 0.0015 , \quad (g_{b\bar{b}}^{(6)})_{\text{exp}} = +0.0962 \pm 0.0063 .$$

While the result for the LH coupling is consistent with the SM prediction, there is a large disagreement between data and SM expectation in the RH sector:

$$(\Delta g_R^{b\bar{b}})_{\text{exp}} = (g_{b\bar{b}}^{(6)})_{\text{exp}} - (g_{b\bar{b}}^{(6)})_{\text{SM}} = (1.9 \pm 0.6) \times 10^{-2} .$$

This deviation could in principle be solved by choosing appropriate couplings for the RH operators in [69]. From the modified RH coupling in (110) we get

$$(\Delta g_R^{b\bar{b}})_{RH} \approx -0.15 \times 10^{-2} \times c_{Z_R}^{\text{eff}}$$

where

$$c_{Z_R}^{\text{eff}} = (c_{R_{Z1}} + 2c_{R_{Z2}}) \frac{(3 \text{ TeV})^2} {\Lambda^2} ,$$

and the numerical value has been obtained assuming $|\tilde{V}_{bb}|^2 \approx 1/2$ (see Sect. 4.3). As can be seen, for $\Lambda = 3$ TeV and $c_{R_{Z1}} = O(1)$, the correction is too small to contribute significantly to the experimental discrepancy in (113). In principle, the effect could be explained assuming $\Lambda = 1$ TeV and $c_{R_{Z1}}, c_{R_{Z2}} = O(1)$. However, as we will show in the following, this possibility is ruled out after taking into account the phenomenological bounds from rare $B$ decays.
7.3 Rare $B$ and $K$ decays: preliminaries

General phenomenological analyses about the role of $Z$-mediated RH currents in $B$ and $K$ decays with a lepton pair in the final state can be found in [44,45]. While most of the results obtained in these two papers can be applied also to the present study, here we go one step forward having determined a series of constraints on the flavour structure of the RH mixing matrix from other processes. We begin this section by introducing the relevant effective Hamiltonians.

For $B_{s,d} \to \mu^+\mu^-$ channels we generalize the SM effective Hamiltonian to

$$H_{\text{eff}} = -\frac{4G_F}{\sqrt{2}} \frac{\alpha}{2\pi s_W} V_{tb}^\dagger V_{tq} \times \left[ Y_{LL}(\bar{b}_L\gamma^\mu s_L)(\bar{\mu}_L\gamma_\mu\mu_L) ight. \\
\left. + Y_{LR}(\bar{b}_L\gamma^\mu s_L)(\bar{\mu}_R\gamma_\mu\mu_R) + Y_{RL}(\bar{b}_R\gamma^\mu s_R)(\bar{\mu}_L\gamma_\mu\mu_L) + Y_{RR}(\bar{b}_R\gamma^\mu s_R)(\bar{\mu}_R\gamma_\mu\mu_R) \right], \quad (116)$$

where $q = d, s$. The overall factor allows for an easy comparison with the SM, where $Y_{LL} - Y_{LR} = Y_0(x_1)$ and $Y_{RR} = Y_{RL} = 0$. The contributions to the $Y$ functions with RH quark currents obtained by means of the effective Lagrangian ($106$) are

$$Y_{RL} - Y_{RR} = -T \frac{\bar{V}_{tb}^* V_{tq}}{V_{tb}^* V_{tq}}, \quad Y_{RL} + Y_{RR} = -(1 - 4s_W^2) T \frac{\bar{V}_{tb}^* V_{tq}}{V_{tb}^* V_{tq}}, \quad (117)$$

where we have defined

$$T = (c_{RZ} + 2c_{RZ}) \frac{4\pi^2 v^2 y_t^2}{g^2 A^2} = 0.55 \times \left( \frac{m_t(m_t)}{163.5 \text{ GeV}} \right)^2 c_{Z_R}^\dagger. \quad (118)$$

Note that $Y_{LL}$ and $Y_{LR}$ are not affected by the effective Lagrangian ($106$).

A simple generalization of the SM effective Hamiltonians can be implemented also to describe $K \to \pi \nu \bar{\nu}$ and $B \to \{X_s, K, K^*\} \nu \bar{\nu}$ decays, with the further simplification that we can neglect operators with $\nu_R$ fields, that we assumed to be heavy. In the $K \to \pi \nu \bar{\nu}$ case we generalize the short-distance effective Hamiltonian to

$$H_{\text{eff}} = \frac{4G_F}{\sqrt{2}} \frac{\alpha}{2\pi s_W} V_{ts}^\dagger V_{td} \times \left[ X_{LL}(K)(\bar{s}_L\gamma^\mu d_L) + X_{RL}(K)(\bar{s}_R\gamma^\mu d_R) \right] \times (\bar{\nu}_L\gamma_\mu\nu_L), \quad (119)$$

where the leading SM top-quark contribution yields $X_{RL} = 0$ and $X_{LL} \equiv X_{SM} = 1.464 \pm 0.041$ [46] (for simplicity we omit the sub-leading charm contribution that will be included in the phenomenological analysis). With these notations the non-standard contribution obtained by means of the effective Lagrangian ($106$) is

$$X_{RL}(K) = -T \frac{\bar{V}_{ts}^* V_{td}}{V_{ts}^* V_{td}}. \quad (120)$$

In the $B \to \{X_s, K, K^*\} \nu \bar{\nu}$ case the general effective Hamiltonian is

$$H_{\text{eff}} = \frac{4G_F}{\sqrt{2}} \frac{\alpha}{2\pi s_W^2} V_{ts}^\dagger V_{ts} \times \left[ X_{LL}(B_s)(\bar{b}_L\gamma^\mu s_L) + X_{RL}(B_s)(\bar{b}_R\gamma^\mu s_R) \right] \times (\bar{\nu}_L\gamma_\mu\nu_L), \quad (121)$$

with

$$X_{RL}(B_s) = -T \frac{\bar{V}_{ts}^* V_{ts}}{V_{ts}^* V_{ts}}. \quad (122)$$
7.4 \(B_{s,d} \rightarrow \mu^+\mu^-\)

When evaluating the amplitude for \(B_s \rightarrow \mu^+\mu^-\) by means of (116) the following simplifications occur

\[
\langle 0 | b \gamma_\mu P_{R,L}\gamma | B^0 \rangle = \pm \frac{1}{2} \langle 0 | b \gamma_\mu \gamma_5 | B^0 \rangle , \quad \langle \bar{\mu}\mu | \bar{\mu}\gamma_\mu P_{R,L}\mu | 0 \rangle = \pm \frac{1}{2} \langle \bar{\mu}\mu | \bar{\mu}\gamma_\mu \gamma_5 | 0 \rangle . \tag{123}
\]

The resulting branching ratio is then obtained from the known SM expression (see e.g. [47]) by making the following replacement

\[
Y_0(x_t) \rightarrow Y_{LL} + Y_{RR} - Y_{RL} - Y_{LR} \equiv Y_{tot} \tag{124}
\]

so that

\[
B(B_s \rightarrow \ell^+\ell^-) = \tau(B_s) \frac{G_F^2}{\pi} \left( \frac{\alpha}{4 \pi s_W^2} \right)^2 F_{B_s}^2 m_{B_s}^2 \sqrt{1 - 4 \frac{m_q^2}{m_{B_s}^2}} |V_{ts}|^2 |Y_{tot}|^2 . \tag{125}
\]

The expression for \(B(B_d \rightarrow \ell^+\ell^-)\) is obtained by replacing \(s\) by \(d\).

Taking into account that \(\hat{V}_{td}^\dagger V_{ts} \approx \hat{s}_{12} e^{i\phi_{21}} / 2\) and \(\hat{V}_{tb}^\dagger V_{td} \approx \hat{s}_{12} e^{i\phi_{23}} / 2\) (see Sect. 4.3 and Sect. 4.3), and using (117), we finally obtain the following expressions for the two branching ratios normalized to the SM:

\[
B(B_s \rightarrow \ell^+\ell^-) = B(B_s \rightarrow \ell^+\ell^-)_{SM} |1 + 7.8 \times \hat{s}_{12} e^{i\phi_{21}} \epsilon_R^{eff}|^2 , \\
B(B_d \rightarrow \ell^+\ell^-) = B(B_d \rightarrow \ell^+\ell^-)_{SM} |1 + 37 \times \hat{c}_{12} e^{i\phi_{23}} \epsilon_R^{eff}|^2 . \tag{126}
\]

The muon channels are those where the experimental searches are closer to the SM predictions. The numerical values of the latter, obtained using the relation of \(B(B_q \rightarrow \mu^+\mu^-)\) to \(\Delta M_q\) pointed out in [47], are

\[
B(B_s \rightarrow \mu^+\mu^-) = (3.2 \pm 0.2) \times 10^{-9} , \quad B(B_d \rightarrow \mu^+\mu^-) = (1.0 \pm 0.1) \times 10^{-10} . \tag{127}
\]

These figures should be compared with the 95% C.L. upper limits from CDF [48] and D0 [49] (in parentheses)

\[
B(B_s \rightarrow \mu^+\mu^-) \leq 3.3 (5.3) \times 10^{-8} , \quad B(B_d \rightarrow \mu^+\mu^-) \leq 1 \times 10^{-8} . \tag{128}
\]

Using the results in (126) these limits imply

\[
|\hat{s}_{12} \epsilon_R^{eff}| < 0.54 , \quad |\hat{c}_{12} \epsilon_R^{eff}| < 0.30 , \tag{129}
\]

where the bounds have been derived taking into account the interference with the SM (and choosing the maximal interference effect). These two limits can be combined to derive the following bound

\[
|\epsilon_R^{eff}| < 0.62 , \tag{130}
\]

which holds independently of any assumption about the value of \(\hat{c}_{12}\). Using this bound in (114) we get

\[
|\langle \Delta g_{RH}^{bb} \rangle| < 1 \times 10^{-3} , \tag{131}
\]

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which, by construction, does not rely on any assumption about the value of \( \tilde{c}_{12} \). It is then clear that, despite the presence of a non-standard coupling of the \( Z \) boson to RH fermions, within our effective theory the constraints from \( \mathcal{B}(B_{d,s} \to \ell^+\ell^-) \) prevent a solution to the \( Z \to b \bar{b} \) anomaly in (113).

The bound (130) has been derived from the experimental bounds on both \( \mathcal{B}(B_s \to \mu^+\mu^-) \) and \( \mathcal{B}(B_d \to \mu^+\mu^-) \) in order to show in a simple manner that there is no room for a sizable contribution to \( Z \to b \bar{b} \) in our framework, independently of the value of \( \tilde{c}_{12} \). As far as the maximal enhancement of \( \mathcal{B}(B_s \to \mu^+\mu^-) \) is concerned, a more stringent bound can be derived taking into account the constraints on the effective \( Z^\mu \bar{b}_R \gamma^\mu s_R \) coupling following from \( \mathcal{B}(B_{s,d} \to X_s\ell^+\ell^-) \) [44,45]. In particular, from the bound reported in [45] one finds

\[
|T| \times \left| \frac{V_{tb}^* V_{ts}}{V_{tb}^* V_{ts}} \right| < 1.07, \tag{132}
\]

at the 90\% C.L. level, or

\[
\left| \tilde{s}_{12} \epsilon_{Z_R}^\text{eff} \right| < 0.15. \tag{133}
\]

Using this result in (126), the maximal enhancement in \( \mathcal{B}(B_s \to \mu^+\mu^-) \) over its SM expectation does not exceed a factor of 5. This should be contrasted to other NP frameworks, in particular to models with non-standard scalar FCNCs, where the present experimental upper bound on \( \mathcal{B}(B_s \to \mu^+\mu^-) \) could easily be saturated.

In case of an \( \mathcal{O}(1) \) deviation from the SM in \( \mathcal{B}(B_s \to \mu^+\mu^-) \), a clear prediction of our framework, following from the analysis of \( \Delta F = 2 \) processes, is the absence of visible deviations from the SM in \( \mathcal{B}(B_d \to \mu^+\mu^-) \). Indeed, as we have seen in Section 6.5 the configuration of the RH matrix that could allow to explain a large \( \mathcal{S}_{\psi\psi} \) asymmetry requires \( \tilde{s}_{12} \approx 1 \) and \( \tilde{c}_{12} < 10^{-2} \). When combined with the bound in (133) this condition implies negligible non-standard effects in \( \mathcal{B}(B_d \to \mu^+\mu^-) \).

**7.5 \( B \to \{X_s, K, K^*\}\nu\bar{\nu} \)**

Following the analysis of Ref. [45], the branching ratios of the \( B \to \{X_s, K, K^*\}\nu\bar{\nu} \) modes in the presence of RH currents can be written as follows

\[
\mathcal{B}(B \to K\nu\bar{\nu}) = \mathcal{B}(B \to K\nu\bar{\nu})_{\text{SM}} \times [1 - 2\eta] \epsilon^2, \tag{134}
\]

\[
\mathcal{B}(B \to K^*\nu\bar{\nu}) = \mathcal{B}(B \to K^*\nu\bar{\nu})_{\text{SM}} \times [1 + 1.31\eta] \epsilon^2, \tag{135}
\]

\[
\mathcal{B}(B \to X_s\nu\bar{\nu}) = \mathcal{B}(B \to X_s\nu\bar{\nu})_{\text{SM}} \times [1 + 0.09\eta] \epsilon^2, \tag{136}
\]

where we have introduced the variables

\[
\epsilon^2 = \frac{|X_{LL}|^2 + |X_{RL}|^2}{|X_{LL}^\text{SM}|^2}, \quad \eta = \frac{-\text{Re}(X_{LL}^* X_{RL})}{|X_{LL}|^2 + |X_{RL}|^2}, \tag{137}
\]

in terms of the Wilson coefficient of the effective Hamiltonian (121).\(^1\) To simplify the notations, here and in the following we omit to specify the meson system in the \( X_{LL,LR} \) functions.

\(^1\) The expressions in Eqs. (134)-(136), as well as the SM figures in (138), refer only to the short-distance contributions to these decays. The latter are obtained from the corresponding total rates subtracting the reducible long-distance effects pointed out in [50].
The updated predictions for the SM branching ratios are \cite{45, 50, 51}
\begin{align}
B(B \to K \nu \bar{\nu})_{SM} &= (3.64 \pm 0.47) \times 10^{-6}, \\
B(B \to K^* \nu \bar{\nu})_{SM} &= (7.2 \pm 1.1) \times 10^{-6}, \\
B(B \to X_s \nu \bar{\nu})_{SM} &= (2.7 \pm 0.2) \times 10^{-5},
\end{align}
to be compared with the experimental bounds \cite{52–54}
\begin{align}
B(B \to K \nu \bar{\nu}) &< 1.4 \times 10^{-5}, \\
B(B \to K^* \nu \bar{\nu}) &< 8.0 \times 10^{-5}, \\
B(B \to X_s \nu \bar{\nu}) &< 6.4 \times 10^{-4}.
\end{align}

The expressions in Eqs. (134)–(136) are valid for wide class of NP model: all models giving rise to the effective Hamiltonian \cite{121}. In our specific framework NP effects are encoded only in the RH sector and the $X_{RL}$ function is given in \cite{122}. The variables $\epsilon$ and $\eta$ then assume the following form:
\begin{align}
\epsilon^2 &= 1 + \frac{T^2}{X_0^2\left(x_t\right)} \left| \frac{\bar{V}_{tb}^* V_{ts}}{V_{tb}^* V_{ts}} \right|^2 \approx 1 + 22.1 \times |\tilde{s}_{12} c_{Z_R}^{\text{eff}}|^2, \\
\eta &= \frac{T}{\epsilon^2 X_0 \left(x_t\right) \text{Re} \left( \frac{\bar{V}_{tb}^* V_{ts}}{V_{tb}^* V_{ts}} \right)} \approx \frac{4.7 \times \tilde{s}_{12} \cos(\phi_{32}^d) c_{Z_R}^{\text{eff}}}{1 + 22.1 \times |\tilde{s}_{12} c_{Z_R}^{\text{eff}}|^2}.
\end{align}

Taking into account the bound on $\tilde{s}_{12} c_{Z_R}^{\text{eff}}$ in Eq. (133), we find that the predictions for the exclusive branching ratios can be enhanced by more than a factor of two over the corresponding SM estimates. On the contrary, the enhancement in the inclusive mode does not exceed 50%. Most important, a clear prediction of RH currents is the anti-correlation of the two exclusive modes: if $B(B \to K \nu \bar{\nu})$ is enhanced then $B(B \to K^* \nu \bar{\nu})$ is suppressed, and vice versa.\footnote{In principle, similar correlations could also be established in the $B \to K \ell^+ \ell^-$ and $B \to K^* \ell^+ \ell^-$ channels. However, in this case the pattern is less clean due to the presence of other effective operators. Moreover, the simplifying assumption of considering only the effective Lagrangian \cite{70} as representative of the dominant RH}

Not surprisingly, the pattern of these three modes is very similar to what observed in Section 4.2. However, in the rare modes the deviations from the SM can in principle be larger than in the charged-current decays.

The correlations among the three $B \to \{X_s, K, K^*\} \nu \bar{\nu}$ modes are in principle affected by the uncertainty on the CP-violating phase $\phi_{32}^d$. However, the same phase enters in $S_{\psi \phi}$. If we require a large $S_{\psi \phi}$ (as hinted by CDF and D0), this uncertainty is strongly reduced, as shown in Figure 2. In this plot we show the expectations of the two exclusive branching ratios for the preferred values of the CP-violating phase $\phi_{32}^d$ as determined from $S_{\psi \phi}$ in Section 6.5. Two points should be noted: 1) only the modulo of $\sin(2\phi_{32}^d)$ enters in the branching ratios of the rare modes (via their $\cos(\phi_{32}^d)$ dependence), as a result, there are only two independent choices corresponding to all the solutions in in Eqs. (100) and Eqs. (103); 2) these two choices give rise to predictions for the $B \to \{K, K^*\} \nu \bar{\nu}$ branching ratios which are almost indistinguishable. As shown in Figure 2, the correlation pattern is very clean and, if observed, would provide a clear confirmation of this framework.
Figure 2: Correlation between $\mathcal{B}(B \to K \nu \bar{\nu})$ and $\mathcal{B}(B \to K^* \nu \bar{\nu})$ in our effective theory. The two bands correspond to the two values of $|\sin(2\phi_3^d)|$ in Eqs. (100), with the uncertainty given by the errors on the SM predictions: blue (dark gray) band for $|\sin(2\phi_3^d)| = 0.95$, orange (light gray) band for $|\sin(2\phi_3^d)| = 0.30$. The black point denotes the SM values with the corresponding error bars.

7.6 $K \to \pi \nu \bar{\nu}$

The SM branching ratios for the two most interesting $K \to \pi \nu \bar{\nu}$ modes can be written as $^{56}59$

$$\mathcal{B}(K^+ \to \pi^+ \nu \bar{\nu}) = \kappa_+ \left( \frac{\text{Im}X_{\text{eff}}}{\lambda^5} \right)^2 + \left( \frac{\text{Re}X_{\text{eff}}}{\lambda^5} - P_c - \delta P_{c,u} \right)^2, \quad (142)$$

$$\mathcal{B}(K_L \to \pi^0 \nu \bar{\nu}) = \kappa_L \left( \frac{\text{Im}X_{\text{eff}}}{\lambda^5} \right)^2, \quad (143)$$

where

$$X_{\text{eff}} = V_{ts}^* V_{td}(X_{LL} + X_{RL})$$

effects is not necessarily a good approximation for these channels. A detailed analysis of $B \to K \ell^+ \ell^-$ and $B \to K^* \ell^+ \ell^-$ in our effective theory goes beyond the purpose of the present paper and we refer to the general model-independent analysis in Ref. $^{55}$. 

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and, as usual, $\lambda = |V_{us}|$, while $\kappa_+ = (5.173 \pm 0.025) \times 10^{-11} (\lambda/0.225)^8$ and $\kappa_L = (2.29 \pm 0.03) \times 10^{-10} (\lambda/0.225)^8$. In the $K^+$ case the dimension-six charm quark corrections and subleading long-distance effects are characterized by $P_c = 0.372 \pm 0.015$ and $\delta P_{c,u} = 0.04 \pm 0.02$, respectively. The $X_{\text{eff}}$ function can be rewritten as

$$X_{\text{eff}} = V^*_{ts}V_{td}X_{\text{SM}}(1 + \xi e^{i\theta})$$

where $X_{\text{SM}} = 1.464 \pm 0.041$ and we have introduced the two real parameters $\xi$ and $\theta$ that vanish in the SM. The explicit expression for these two parameters in our framework is

$$\xi e^{i\theta} = -\frac{T}{X_{\text{SM}}} \frac{\bar{V}_{ts}V_{td}}{V_{ts}^*V_{td}} \approx 5.6 \times 10^2 \times \bar{c}_{12} \bar{s}_{12} e^{i(\phi_{d1} + \beta - \beta_s)} c_{\text{eff}} Z_R.$$  (146)

If we ignore the constraints from the $\Delta F = 2$ processes there is certainly a large room for non-standard effects in $K \to \pi\nu\bar{\nu}$ decays, even taking into account the bound in Eq. (133). The situation changes if we implement the constraints on $\bar{c}_{12}$ derived from $\varepsilon_K$, under the hypothesis of a large non-standard contribution to $S_{\psi\phi}$, analysed in Section 6.5. Here we should distinguish two cases:

- If we can neglect the contribution of $Q_{d1}^{LR}$ to $\Delta F = 2$ amplitudes, namely if we need to implement the bounds in Eqs. (102), then there is room for $O(1)$ deviations from the SM predictions in both rare modes for a sizable range of the phase $\phi_{d1}$. This is shown by the red (dark) points in Fig. 3 which are obtained imposing this constraint. As can be seen, larger deviations from the SM are also possible, but in this case the phase $\phi_{d1}$ has to be tuned such that $\sin(2\phi_{d1}) \approx 0$, in order to avoid the $\varepsilon_K$ constraint. This fined-tuned configuration gives rise to the two bands of red (dark) points in Fig. 3 with large enhancements of only one of the two branching ratios. As noted in [62] this structure is characteristic of all NP frameworks where the phase in $\Delta S = 2$ amplitudes is the square of the CP-violating phase in $\Delta S = 1$ FCNC amplitudes (this is for instance what happens in the Little Higgs model with $T$ parity [63]).

- If the contribution of $Q_{d1}^{LR}$ to $\Delta F = 2$ amplitudes is dominant, namely if we need to implement the stringent bounds in Eqs. (104), then the situation is more constrained. Here we can expect a visible deviation from the SM only if the phase $\phi_{d1}$ is tuned such that $\sin(\phi_{d1} - \beta + \beta_s) \approx 0$, where the bounds (104) become less effective. This give rise to the narrow band of green (light) points in Fig. 3. This correlation is very different from the one pointed out above, since in this case the leading CP-violating phase in $\Delta S = 2$ amplitudes is not the square of the CP-violating phase appearing in $\Delta S = 1$ FCNC amplitudes.

If we relax the assumption of sizable NP contributions to $S_{\psi\phi}$ the predictions of these two modes do not change substantially: the characteristic structures in Fig. 3 are indeed only due to the $\varepsilon_K$ constraint. As anticipated, the situation may change only if we could ignore the constraints from the $\Delta S = 2$ processes, assuming the corresponding Wilson coefficients are accidentally suppressed. However, we consider this situation highly fine-tuned. Actually it should be stressed that the maximal enhancements shown in Fig. 3 also require a considerable amount of fine-tuning, both on the phase $\phi_{d1}$ and on the ratio of $\Delta S = 2$ over $\Delta S = 1$ Wilson coefficients.
Figure 3: Correlations between $B(K^+ \to \pi^+ \nu \bar{\nu})$ and $B(K_L \to \pi^0 \nu \bar{\nu})$ in our effective theory, taking into account the $\varepsilon_K$ constraint. The red (dark) and green (light) points are obtained imposing the $\varepsilon_K$ constraint under the assumption of negligible or dominant contribution from $Q_1^{LR}$, respectively. The dashed line is the Grossman-Nir bound [60]. The vertical band correspond to the experimental result in [61] (1 $\sigma$ range) and the black cross to the SM prediction.

7.7 Summary

The main results of the $\Delta F = 1$ analysis can be summarized as follows:

- The constraints from $B_{s,d} \to \mu^+ \mu^-$ eliminate the possibility of removing the known anomaly in the $Z \to b\bar{b}$ decay with the help of right-handed currents.
- Contributions from RH currents to $B_{s,d} \to \mu^+ \mu^-$, $B \to \{X_s, K, K^*\} \nu \bar{\nu}$, and $K \to \pi \nu \bar{\nu}$ can all be significant, although this is not a general prediction of the model. If the deviations from the SM are sizable, the effects exhibit interesting patterns of correlations.
- In the $B_s \to \mu^+ \mu^-$ case the branching ratio can receive an $O(1)$ enhancement over its SM expectation, but it cannot get close to its present experimental bound due to the constraint from $B \to X_s l^+ l^-$. If the RH contribution to $S_{\psi\phi}$ is large, no significant enhancement is expected in $B_d \to \mu^+ \mu^-$. 
- If the RH contribution to $S_{\psi\phi}$ is large, the pattern of possible enhancement/suppression in $B \to \{K, K^*\} \nu \bar{\nu}$ is unambiguous, as shown in Fig. 2.
- The pattern of possible enhancement/suppression in the to $K \rightarrow \pi \nu \bar{\nu}$ modes is largely independent of possible RH contributions to $S_{\psi\phi}$, but it could help to disentangle the case where non-standard $\Delta F = 2$ amplitudes are dominated by RR or RL operators.

8 Comparison with MFV and explicit LR models

8.1 RH currents vs. MFV: general considerations

The two effective theories are apparently very similar: the low-energy particle content is the same and in both cases we have a flavour group broken only by two Yukawas. However, the flavour groups are different: $SU(3)_Q \times SU(3)_U \times SU(3)_D$ in the MFV case [14] vs. $SU(3)_L \times SU(3)_R$ in the present framework. The larger flavour symmetry of the MFV set-up implies a much more constrained structure. In particular:

- If the normalization of the two Yukawa couplings is the same as in the SM, as expected with a single Higgs doublet, we are in the so-called CMFV regime [30]. In this case deviations from the SM are small in all observables. In particular, there is no hope to explain a large $S_{\psi\phi}$ asymmetry, contrary to what happens with RH currents.

- The expectation of a vanishingly small $S_{\psi\phi}$ in the MFV case remains true also with two Higgs doublets and large value of $\tan \beta$ if we assume that the Yukawa couplings are the only sources of $CP$ violation, as originally assumed in [14]. In this case the only large deviation from the SM is a potentially large $B_s \rightarrow \mu^+\mu^-$ rate that could easily be just below the present experimental bound. As we have seen, in the RH framework $B_s \rightarrow \mu^+\mu^-$ could be enhanced over its SM expectation, but the enhancement cannot be as large as in the MFV case at large $\tan \beta$ because of the $B_s \rightarrow X_s\ell^+\ell^-$ constraint.

8.2 Comparison with MFV with flavour-blind phases

The phenomenology of MFV models can be quite different from the case discussed above if one relaxes the assumption that the Yukawa couplings are the only sources of $CP$ violation, or in the GMFV framework, as denoted in [64]. In particular, it has been recently shown that in a two Higgs doublet model (2HDM) with MFV, large $\tan \beta$, and flavour-blind $CP$-violating phases, it is possible to generate a large $S_{\psi\phi}$ asymmetry and, as a consequence, automatically soften the anomalies in $S_{\psi K_S}$ and $\varepsilon_K$ in a correlated manner [41]. In this set-up the NP contributions responsible for a large $S_{\psi\phi}$ are due to the exchange of heavy neutral scalars. It is then interesting to compare this solution to the $S_{\psi\phi}$ problem with the one considered here, where NP contributions originate presumably from new heavy gauge bosons.

Concentrating first on the high energy scales, at which new particles are integrated out let us emphasize that the relevant scales of right-handed currents are by roughly an order of magnitude larger than the allowed masses of neutral scalars. Also the Lorentz structure of the operators generated at the high scale is different in these models. As seen in Eqs. (79) and (95), the presence of RH currents selects from the list in (71) non-vanishing initial conditions for $Q_1^{VRR}$ and $Q_1^{L_R}$, while the leading $SU(2)_L \times U(1)_Y$ invariant operator induced by Higgs exchanges at the high scale is $Q_2^{LR}$. As a consequence of this different operator structure, the NP contributions from RH currents are governed by the RG parameters $P_1^{VRR} = P_1^{VLL}$.
and $P^{LR}_1$, while the ones in the 2HDM by $P^{LR}_2$. It is then interesting to observe that model independently $P^{LR}_1$ and $P^{LR}_2$ are roughly of the same magnitude but opposite sign. While this sign difference is not relevant in view of unknown signs of the coefficients involved in a model with RH currents, it could play some role when a concrete model with RH currents is analyzed.

In the case of the 2HDM with MFV the pattern of NP contributions to $K^0 - \bar{K}^0$, $B^0_d - \bar{B}^0_d$ and $B^0_s - \bar{B}^0_s$ mixings is governed by external quark masses implying naturally largest effects in the $B^0_s - \bar{B}^0_s$, followed by an order of magnitude $O(m_d/m_s)$ smaller effects in $B^0_d - \bar{B}^0_d$ mixing and negligible effects in $\varepsilon_K$. Still the impact of NP on the $B^0_d - \bar{B}^0_d$ mixing, uniquely following from the requirement of fitting a large $S_{\psi\phi}$, softens automatically the $\varepsilon_K$ anomaly through the increase of the true value of $\sin 2\beta$ resulting from the fit to $S_{\psi K_S}$. It should also be noted that the increase of $\sin 2\beta$ in the 2HDM automatically favors larger values of $|V_{ub}|$ than those found in $B \to \pi e\nu$, but the model does not offer an explanation why the exclusive and inclusive determinations give different values of $|V_{ub}|$.

As we have shown in the previous sections, the pattern of deviations from the SM is significantly different in the case of RH currents, although some points are common to both frameworks. In particular:

- RH currents provide a natural explanation of the different values of $|V_{ub}|$ following from inclusive and exclusive decays, selecting the inclusive determination as the one giving the true value of $|V_{ub}|$. This value is significantly higher than the corresponding value obtained from the SM fits and, when combined with the tree level measurement of the angle $\gamma$, implies $\sin 2\beta = 0.77 \pm 0.05$.

- Similarly to the MFV-2HDM case, the modified determination of $\sin 2\beta$ removes the $\varepsilon_K$ anomaly even in the absence of direct NP contributions to this observable.

- In contrast to the MFV-2HDM case, the RH current contributions to $B^0_d - \bar{B}^0_d$ mixing are constrained to be negligible. As a result we should expect $S_{\psi K_S} = \sin 2\beta$, implying a value of $S_{\psi K_S}$ significantly larger than what determined from experiments.

- The model has sufficient number of parameters that it can naturally generate a large value of $S_{\psi\phi}$ without any conflict with other data. Contrary to the MFV-2HDM case, the large $S_{\psi\phi}$ does not imply $B_s \to \ell^+\ell^-$ close to the present experimental limit. Moreover $B_d \to \ell^+\ell^-$ receives only small NP contribution so that the MFV relation between the the branching ratios for $B_s \to \ell^+\ell^-$ and $B_d \to \ell^+\ell^-$ can be strongly violated.

We conclude then that both models provide interesting solutions to the existing anomalies but:

- the 2HDM with MFV cannot provide the explanation of different values of $|V_{ub}|$ following from inclusive and exclusive semi-leptonic decays;

- the model with RH currents provides this explanation but this feature combined with tiny contributions to $B^0_d - \bar{B}^0_d$ mixing implies a $\approx 2\sigma$ discrepancy between the predicted value of $S_{\psi K_S}$ and its experimental determination.

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8.3 Comparison with explicit left-right models.

As anticipated in the introduction, there exist several analyses of flavour observables in explicit left-right models (see e.g. [6, 7] and references therein). In all these papers flavour observables are a key ingredient to determine the bounds on the masses of the massive RH gauge bosons. Most of the existing analyses are focused on the minimal version of the model [2–4], that is characterized by the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, by a discrete symmetry connecting the two $SU(2)$ groups, and by the minimal choice for the Higgs sector necessary to achieve the two-step breaking $SU(2)_L \times SU(2)_R \times U(1)_{B-L} \rightarrow SU(2)_L \times U(1)_Y \rightarrow U(1)_{\text{em}}$. A more general analysis where the latter two hypotheses are relaxed can be found in [65].

While some features of our effective theory approach can be applied also to these explicit models, there are a few important distinctive features of our analysis, related to the assumptions behind the Yukawa interactions.

- In our approach we make no assumptions about the existence of a $L \leftrightarrow R$ discrete symmetry under which the Yukawa interactions must be invariant. This symmetry, that is usually enforced in the minimal models, force the Yukawa matrices to be exactly (or approximately) LR symmetric [7]. As a result, $\tilde{V}$ should have a hierarchical structure identical (or very similar) to the one of the CKM matrix. While this is a nice feature as far as protecting FCNCs, it would prevent us to solve the $V_{ub}$ problem, which was one of the main motivation of our analysis. As far as the structure of $\tilde{V}$ and its impact in charged-currents are concerned, our analysis is thus more general than existing analysis in explicit LR models.

- While we do not impose any $L \leftrightarrow R$ discrete symmetry, we assume a minimal structure for the breaking of the flavour group. We assume only two independent Yukawa couplings, with an extra protective symmetry forcing them to act only in the up- and down-type sector in the dimension-four operators, as shown in Eq. (7). This extra assumptions have been imposed to have an efficient suppression of FCNCs in the left-handed sector, and to avoid scalar FCNCs at the tree level. While these assumptions are quite reasonable and can easily be implemented in explicit LR models, they do not represent the most general possibility. It is worth to stress that some of our phenomenological conclusions, such as the absence of NP effects in $B_d$ mixing, after we require large NP effects in $B_s$ mixing, do depend on this assumption.

9 Conclusions

The possibility that at very short distance scales the nature is left-right symmetric appears to be intriguing. As at low energy scales the parity is maximally broken and charged weak interactions exhibit left-handed structure, the right-handed weak currents, if present in nature, must be coupled to new heavy gauge bosons that are at least by one order of magnitude heavier than the $W^\pm$ and $Z$ bosons of the SM.

While such heavy gauge bosons could be discovered at the LHC in the coming years, they can also manifest themselves in low energy processes. In our model they are represented by new effective operators containing right-handed currents with their Wilson coefficients encoding the information about the fundamental theory, in particular the relevant couplings.
In the present paper we have analyzed the impact of right-handed currents in both charged- and neutral-current flavour-violating processes by means of an effective theory approach. To this end we have assumed a left-right symmetric flavour group, commuting with an underlying $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ global symmetry, broken only by two Yukawa couplings. Having identified the leading six dimension operators in this model, we performed a rather detailed analysis of those low energy observables, which could help to support or falsify the presence of this NP at scales probed soon directly by the LHC.

The central role in our model is played by a new unitary matrix $\tilde{V}$ that controls flavour-mixing in the right-handed sector. Using the data on the tree level charged current transitions $u \to d$, $u \to s$, $b \to u$ and $b \to c$ and the unitarity of $\tilde{V}$ we could determine the structure of this matrix and demonstrate, following other authors, that the tension between inclusive and exclusive determinations of $|V_{ub}|$ can be solved with the help of right-handed currents. The resulting true value of $|V_{ub}|$ turns out to be $(4.1 \pm 0.2) \times 10^{-3}$, in the ball park of inclusive determinations. The novel feature of our analysis as compared with [12] is the determination of the full right-handed matrix, and not only its selected elements, making use of unitarity. We also find that while RH currents are very welcome to solve the “$|V_{ub}|$ problem” they do not have a significant impact on the determination of $|V_{cb}|$ (as also pointed out in [15]).

Having determined the size and the flavour structure of right-handed currents that is consistent with the present data on tree level processes and which removes the “$|V_{ub}|$-problem”, we have investigated how this NP would manifest itself in neutral current processes, including particle-antiparticle mixing, $Z \to bb$, $B_{s,d} \to \mu^+\mu^-$, $B \to \{X_s,K,K^*\}\nu\bar{\nu}$ and $K \to \pi\nu\bar{\nu}$ decays. We have also addressed the possibility to explain a non-standard CP-violating phase in $B_s$ mixing in this context, and made the comparison with other predictive new-physics frameworks addressing the same problem.

The main messages from this analysis are as follows:

- The presence of RH currents in the model, in conjunction with the already present SM left-handed currents generates in addition to the operators with $(V + A) \times (V + A)$ Dirac structure, also left-right operators $(V - A) \times (V + A)$. The contributions of the latter are known to be strongly enhanced at low energies through renormalization group effects and in the case of $\varepsilon_K$ and $\Delta M_K$ through chirally enhanced hadronic matrix elements of $(V - A) \times (V + A)$ operators. Consequently these observables put severe constraints on the model parameters (as also known from various studies in explicit LR models [6]).

- The desire to generate large CP-violating effects in $B_s$-mixing, hinted for by the enhanced value of $S_{\psi\phi}$ observed by the CDF and D0 collaborations, in conjunction with the $\varepsilon_K$-constraint, implies additional constraints on the shape of $\tilde{V}$. In particular $\delta_{12} \ll 1$ and consequently $\delta_{12} \approx 1$. The pattern of deviations from the SM in this model is then as follows.

- The $S_{\psi\phi}$ and $\varepsilon_K$ anomalies can be understood.

- As a consequence of the large value of $\delta_{12}$, it should be possible to resolve the presence of RH currents also in $s \to u$ charged-current transitions. Here RH currents imply a $\mathcal{O}(10^{-3})$ deviation in the determination of $|V_{us}|$ from $K \to \pi\ell\nu$ and $K \to \ell\nu$ decays.

- The “true value” of $\sin 2\beta$ determined in our framework, namely the determination of the CKM phase $\beta$ on the basis of the tree-level processes only, and in particular of $|V_{ub}|$, is $\sin 2\beta = 0.77 \pm 0.05$. This result is roughly $2\sigma$ larger than the measured value
\( S_{\psi K_S} = 0.672 \pm 0.023 \). This is a property of any explanation of the “\(|V_{ub}|\)-problem” by means of RH currents, unless the value of \(|V_{ub}|\) from inclusive decays will turn out to be much lower than determined presently. In general, such discrepancy could be solved by a negative new CP-violating phase in \( B_d^0 - \bar{B}_d^0 \) mixing. However, we have demonstrated that this is not possible in the present framework once the \( \varepsilon_K \) constraint is imposed and large \( S_{\psi\phi} \) is required. Thus we point out that simultaneous explanation of the “\(|V_{ub}|\)-problem” and of \( S_{\psi K_S} = 0.672 \pm 0.023 \) is problematic through RH currents alone.

- The present constraints from \( B_{s,d} \to \mu^+\mu^- \) eliminate the possibility of removing the known anomaly in the \( Z \to b\bar{b} \) decay with the help of right-handed currents. On top of it, the constraint from \( B \to X_s l^+l^- \) precludes \( B_s \to \mu^+\mu^- \) to be close to its present experimental bound. Moreover NP effects in \( B_d \to \ell^+\ell^- \) are found generally smaller than in \( B_s \to \ell^+\ell^- \).

- Contributions from RH currents to \( B \to \{X_s, K, K^*\}\nu\bar{\nu} \) and \( K \to \pi\nu\bar{\nu} \) decays can still be significant. Most important, the deviations from the SM in these decays would exhibit a well-defined pattern of correlations.

We have compared this NP scenario with the general MFV framework and with more explicit NP models. Particularly interesting is the comparison with the 2HDM with MFV, large \( \tan \beta \), and flavour-blind CP-violating phases, where the \( S_{\psi\phi} \) and \( \varepsilon_K \) anomalies can also be accommodated \[41\]. What clearly distinguishes these two models at low-energies is how they face the “\(|V_{ub}|\)-problem” (which can be solved only in the RH case) and the “\( \sin 2\beta - S_{\psi K} \)” tension” (which can be softened only in the 2HDM case). But also the future results on rare \( B \) and \( K \) decays listed above could in principle help to distinguish these two general NP frameworks.

Restricting the discussion to these two NP frameworks, it appears that a model with an extended scalar sector and right-handed currents could provide solutions to all the existing tensions in flavour physics simultaneously. This possibility can certainly be realized in explicit left-right symmetric models, where an extended Higgs sector is also required to break the extended gauge symmetry. However, these extensions contain many free parameters and clear cut conclusions on the pattern of flavour violation cannot be as easily reached as it was possible in the simple framework considered here and in \[41\].

We are looking forward to the upcoming experiments at the LHC, future \( B \) factories, and rare \( K \) decay experiments, that should be able to shed more light on the role of right-handed currents in flavour physics.

Acknowledgments

We thank Wolfgang Altmannshofer and Tillmann Heidsieck for useful discussions. G.I. would like to thank Jernej Kamenik and Martin Gonzalez-Alonso for useful discussions in the early stage of this work. We all thank the Galileo Galilei Institute for Theoretical Physics for the hospitality and partial support during the completion of this work. This research was partially supported by the Cluster of Excellence ‘Origin and Structure of the Universe’, by the Graduiertenkolleg GRK 1054 of DFG, by the German ‘Bundesministerium für Bildung und Forschung’ under contract 05H09W0E, and by the EU Marie Curie Research Training Network contract MTRN-CT-2006-035482 (Flavianet).
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