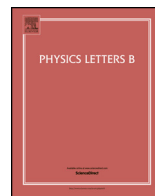


Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletbA unified leptoquark model confronted with lepton non-universality in B -meson decaysT. Faber^a, M. Hudec^b, M. Malinský^{b,*}, P. Meinzinger^a, W. Porod^a, F. Staub^{c,d}^a Institut für Theoretische Physik und Astrophysik, Uni Würzburg, Germany^b Institute for Particle and Nuclear Physics, Charles University, Czech Republic^c Institute for Theoretical Physics (ITP), Karlsruhe Institute of Technology, Engesserstraße 7, D-76128 Karlsruhe, Germany^d Institute for Nuclear Physics (IKP), Karlsruhe Institute of Technology, Hermann-von-Helmholtz-Platz 1, D-76344 Eggenstein-Leopoldshafen, Germany

ARTICLE INFO

Article history:

Received 7 September 2018

Received in revised form 16 October 2018

Accepted 25 October 2018

Available online 30 October 2018

Editor: A. Ringwald

ABSTRACT

The anomalies in the B -meson sector, in particular $R_{K^{(*)}}$ and $R_{D^{(*)}}$, are often interpreted as hints for physics beyond the Standard Model. To this end, leptoquarks or a heavy Z' represent the most popular SM extensions which can explain the observations. However, adding these fields by hand is not very satisfactory as it does not address the big questions like a possible embedding into a unified gauge theory. On the other hand, light leptoquarks within a unified framework are challenging due to additional constraints such as lepton flavor violation. The existing accounts typically deal with this issue by providing estimates on the relevant couplings. In this letter we consider a complete model based on the $SU(4)_C \otimes SU(2)_L \otimes U(1)_R$ gauge symmetry, a subgroup of $SO(10)$, featuring both scalar and vector leptoquarks. We demonstrate that this setup has, in principle, all the potential to accommodate $R_{K^{(*)}}$ and $R_{D^{(*)}}$ while respecting bounds from other sectors usually checked in this context. However, it turns out that $K_L \rightarrow e^\pm \mu^\mp$ severely constraints not only the vector but also the scalar leptoquarks and, consequently, also the room for any sizeable deviations of $R_{K^{(*)}}$ from 1. We briefly comment on the options for extending the model in order to conform this constraint. Moreover, we present a simple criterion for all-orders proton stability within this class of models.

© 2018 The Authors. Published by Elsevier B.V. This is an open access article under the CC BY license (<http://creativecommons.org/licenses/by/4.0/>). Funded by SCOAP³.

1. Introduction

In recent years a few anomalies in the B -meson sector have been observed by different experiments. The most striking one is a $3.5\text{-}\sigma$ deviation in the ratios

$$R_{D^{(*)}} = \frac{\Gamma(\bar{B} \rightarrow D^{(*)} \tau \bar{\nu})}{\Gamma(\bar{B} \rightarrow D^{(*)} l \bar{\nu})} \quad (l = e, \mu),$$

with

$$R_D^{\text{exp}} = 0.388 \pm 0.047, \quad R_{D^*}^{\text{exp}} = 0.321 \pm 0.021, \quad (1)$$

from the Standard Model (SM) lepton universality expectations

$$R_D^{\text{SM}} = 0.300 \pm 0.010, \quad R_{D^*}^{\text{SM}} = 0.252 \pm 0.005. \quad (2)$$

This was first reported by BaBar [1,2] consistent with measurements by Belle [3–5]. Recently this has been confirmed by LHCb in case of R_{D^*} [6] at the $2.1\text{-}\sigma$ level. Additional deviations from lepton universality have recently been reported by LHCb in the ratio

$$R_K = \frac{\Gamma(\bar{B} \rightarrow \bar{K} \mu^+ \mu^-)}{\Gamma(\bar{B} \rightarrow \bar{K} e^+ e^-)} = 0.745_{-0.074}^{+0.090} \pm 0.036, \quad (3)$$

$$R_{K^*} = \frac{\Gamma(\bar{B} \rightarrow \bar{K}^* \mu^+ \mu^-)}{\Gamma(\bar{B} \rightarrow \bar{K}^* e^+ e^-)} = 0.69_{-0.07}^{+0.11} \pm 0.05 \quad (4)$$

in the dilepton invariant mass bin $1 \text{ GeV}^2 \leq q^2 \leq 6 \text{ GeV}^2$ [7,8]. These ratios are predicted to be 1 within the SM and are practically free from theoretical uncertainties. Equally intriguing is a discrepancy in the angular observables in the rare $\bar{B} \rightarrow \bar{K} \mu^+ \mu^-$ decays measured by LHCb [9] which, however, is subject to significant hadronic uncertainties [10,11]. While the individual discrepancies are between 2 and 3 σ , they all point in the same direction and

* Corresponding author.

E-mail addresses: thomas.faber@physik.uni-wuerzburg.de (T. Faber), hudec@ipnp.troja.mff.cuni.cz (M. Hudec), malinsky@ipnp.troja.mff.cuni.cz (M. Malinský), peter.meinzinger@stud-mail.uni-wuerzburg.de (P. Meinzinger), porod@physik.uni-wuerzburg.de (W. Porod), florian.staub@kit.edu (F. Staub).

amount to more than $4.5\text{-}\sigma$ deviations once combined in a fit [12, 13].

In Refs. [14,15] it has been shown that the deviations in R_D and R_K can be explained by an effective model adding one generation of scalar leptoquarks (LQs) with the quantum numbers of the right-handed d -quark and an additional scalar gauge singlet which couples to the LQs. However, it has been shown that this leads to a too large rate for $b \rightarrow s\nu\nu$ [16]. In Ref. [17] another model with two different LQs, one with gauge quantum numbers of the right-handed d -quark and one with charge $4/3$, has been presented which explains also neutrino masses at the 2-loop level. As has been shown in Refs. [18–25], another possibility to successfully accommodate the data is to use vector LQs. A somewhat more complete model containing two types of vector LQs to explain the two-photon excess, based on a Froggatt–Nielsen ansatz for the required coupling structures, has been presented in [26]. Beside the above mentioned violations of lepton-universality this model is also compatible with the neutrino data. Another possibility is that the required leptoquarks are bound states of strongly interacting fermions [27].

Most of these settings are effective models containing just the pieces required to account for the discussed experimental observations, which is clearly the first logical step to make when a new signal shows up. However, eventually one would like to understand the observations from a more fundamental perspective. Several attempts in this direction exist already in the literature [27–38]. Of course, the most attractive scenario would be a UV completion compatible with theoretical requirements like gauge-coupling unification with the potential to explain also the observed dark matter relic density.

From the GUT perspective the Pati–Salam (PS) model [39] emerges as the first and very natural candidate for a low-energy gauge framework featuring vector as well as scalar leptoquarks within a simple dynamical and renormalizable scheme. However, the Kibble–Zurek mechanism of the early-Universe monopole creation [40] suggests that the PS-breaking should occur above the inflation scale [41]. It is therefore advisable to choose $SU(4)_C \otimes SU(2)_L \otimes U(1)_R$ instead as a gauge group of a potentially viable model (as in Refs. [42,43]) which, indeed, does not suffer from the monopole issue.

The structure of this letter is as follows: in Section 2 we present the model and discuss the possibilities to obtain leptoquarks with masses in the TeV range. In Section 3 we discuss in which parts of the parameter space the B anomalies could be accounted for and what are the constraints from the existing low energy data. In Section 4 we draw our conclusions. In Appendix A we demonstrate that in this class of models proton remains stable to all orders in perturbation theory.

2. Model description

In what follows we consider the model proposed by Fileviez-Perez and Wise in Ref. [43]. For convenience, we briefly outline it here, focusing on the features related to flavor physics.

The model is based on the gauge group $G = SU(4)_C \otimes SU(2)_L \otimes U(1)_R$, where the first factor unifies the three colors of quarks with the lepton number. This group is spontaneously broken to $G_{SM} = SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ and further down to $G_{vac} = SU(3)_C \otimes U(1)_Q$, following the branching rules

$$Y = R + 1/2[B - L], \quad Q = T_L^3 + Y, \quad (5)$$

where¹

$$[B - L] = \sqrt{8/3} T_C^{15} = \text{diag} \left(+\frac{1}{3}, +\frac{1}{3}, +\frac{1}{3}, -1 \right). \quad (6)$$

The matching condition for the QCD coupling at the scale, where $SU(4)_C$ is broken, is simply $g_3 = g_4$.

The entire field content of the model is summarised in Table 1. We also include information about other $U(1)$ charges which we need in Appendix A where the details of the baryon number conservation and lepton number violation are discussed.

The SM fermions together with the right-handed neutrinos are combined into three quadruplets under $SU(4)_C$ appearing in three copies representing different generations. On top of that, three fermionic gauge singlets N necessary to generate the correct neutrino masses via inverse seesaw [44] are added.

The gauge field sector corresponding to $SU(4)_C$ consists of the gluons, Z' and a vector leptoquark $X \sim (3, 1, +\frac{2}{3})$ which mediates flavor violating processes such as $K_L \rightarrow e^\pm \mu^\mp$. This tight constraint implies that, for standard-size couplings, the mass of the vector leptoquark has to be at least of the order of 1.6×10^3 TeV [45].²

The scalar sector consists of three multiplets χ , Φ and H , see Table 1. The most general renormalizable scalar potential for these fields reads

$$\begin{aligned} V = & m_H^2 |H|^2 + m_\chi^2 |\chi|^2 + m_\Phi^2 \text{Tr}(|\Phi|^2) + \lambda_1 |H|^2 |\chi|^2 \\ & + \lambda_2 |H|^2 \text{Tr}(|\Phi|^2) + \lambda_3 |\chi|^2 \text{Tr}(|\Phi|^2) + (\lambda_4 H_i^\dagger \chi^\dagger \Phi^i \chi + \text{h.c.}) \\ & + \lambda_5 H_i^\dagger \text{Tr}(\Phi_j^\dagger \Phi^i) H^j + \lambda_6 \chi^\dagger \Phi^i \Phi_j^\dagger \chi + \lambda_7 |H|^4 + \lambda_8 |\chi|^4 \\ & + \lambda_9 \text{Tr}(|\Phi|^4) + \lambda_{10} (\text{Tr}|\Phi|^2)^2 + \left(\lambda_{11} H_i^\dagger \text{Tr}(\Phi^i \Phi^j) H_j^\dagger \right. \\ & + \lambda_{12} H_i^\dagger \text{Tr}(\Phi^i \Phi^j \Phi_j^\dagger) + \lambda_{13} H_i^\dagger \text{Tr}(\Phi^i \Phi_j^\dagger \Phi^j) + \text{h.c.} \left. \right) \\ & + \lambda_{14} \chi^\dagger |\Phi|^2 \chi + \lambda_{15} \text{Tr}(\Phi_i^\dagger \Phi^j \Phi_j^\dagger \Phi^i) \\ & + \lambda_{16} \text{Tr}(\Phi_i^\dagger \Phi^j) \text{Tr}(\Phi_j^\dagger \Phi^i) + \lambda_{17} \text{Tr}(\Phi_i^\dagger \Phi_j^\dagger) \text{Tr}(\Phi^i \Phi^j) \\ & + \lambda_{18} \text{Tr}(\Phi_i^\dagger \Phi_j^\dagger \Phi^i \Phi^j) + \lambda_{19} \text{Tr}(\Phi_i^\dagger \Phi_j^\dagger \Phi^j \Phi^i) \end{aligned} \quad (7)$$

with $|H|^2 = H_i^\dagger H^i$, $|\chi|^2 = \chi^\dagger \chi$, $|\Phi|^2 = \Phi_i^\dagger \Phi^i$ where i, j are the $SU(2)$ indices. The trace is taken over the $SU(4)_C$ indices only. Notice that the terms proportional to $\lambda_{11}, \dots, \lambda_{19}$ have been omitted in the original paper [43]; we include them here for completeness.

The breaking of the $SU(4)_C$ group as well as the electroweak symmetry breaking is triggered by the corresponding vacuum expectation values³ (VEVs)

$$\begin{aligned} \langle \chi \rangle &= \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_\chi \end{pmatrix}, & \langle H \rangle &= \frac{1}{\sqrt{2}} \begin{bmatrix} 0 \\ v_1 \end{bmatrix}, \\ \langle \Phi \rangle &= \frac{1}{2\sqrt{6}} \begin{pmatrix} \mathbb{1} & 0 \\ 0 & -3 \end{pmatrix} \otimes \begin{bmatrix} 0 \\ v_2 \end{bmatrix}, \end{aligned} \quad (8)$$

which are parametrised by $v_1 = v_{ew} \sin \beta$, $v_2 = v_{ew} \cos \beta$, with $v_{ew} \simeq 246$ GeV. The SM-like Higgs h is a superposition of the fields $\text{Re}(\{H_1^0, H_2^0, \chi^0\})$.

The fermion masses are generated by the following interactions between the scalars and fermions:

$$\begin{aligned} -\mathcal{L}_Y = & f_u^c Y_1 H F_L + f_u^c Y_2 \Phi F_L + f_d^c Y_3 H^\dagger F_L \\ & + f_d^c Y_4 \Phi^\dagger F_L + f_u^c Y_5 \chi N + \frac{1}{2} N \mu N + \text{h.c.} \end{aligned} \quad (9)$$

² Note that this limit may be significantly reduced if the freedom in the mixing in the charged leptoquark currents is fully exploited, see, e.g., [46] and references therein.

³ The round and square brackets are used to distinguish between the $SU(4)_C$ and $SU(2)_L$ multiplets, respectively.

¹ We use the square brackets here in order to indicate an indivisible symbol.

Table 1

The field content of the model together with all gauge quantum numbers for the different regimes as well as charges under several other global $U(1)$'s. Whenever $L = L'$ only the common value is displayed.

	G	G_{SM}	G_{vac}	$[B - L]$	F	M	B	L, L'
Fermions								
$F_L = \begin{pmatrix} Q \\ L \end{pmatrix}$	$(4, 2, 0)$	$Q \ (3, 2, 1/6)$ $L \ (1, 2, -1/2)$	$u \ (3, 2/3)$ $d \ (3, -1/3)$ $\nu \ (1, 0)$ $e \ (1, -1)$	$\begin{pmatrix} +1/3 \\ -1 \end{pmatrix}$	+1	+1	$\begin{pmatrix} +1/3 \\ 0 \end{pmatrix}$	$\begin{pmatrix} 0 \\ +1 \end{pmatrix}$
$f_u^c = (u^c \ \nu^c)$	$(\bar{4}, 1, -1/2)$	$u^c \ (\bar{3}, 1, -2/3)$ $\nu^c \ (1, 1, 0)$	$(\bar{3}, -2/3)$ $(1, 0)$	$\begin{pmatrix} -1/3 & 1 \end{pmatrix}$	-1	-1	$\begin{pmatrix} -1/3 & 0 \end{pmatrix}$	$\begin{pmatrix} 0 & -1 \end{pmatrix}$
$f_d^c = (d^c \ e^c)$	$(\bar{4}, 1, 1/2)$	$d^c \ (\bar{3}, 1, 1/3)$ $e^c \ (1, 1, 1)$	$(\bar{3}, 1/3)$ $(1, 1)$	$\begin{pmatrix} -1/3 & 1 \end{pmatrix}$	-1	-1	$\begin{pmatrix} -1/3 & 0 \end{pmatrix}$	$\begin{pmatrix} 0 & -1 \end{pmatrix}$
N	$(1, 1, 0)$	$(1, 1, 0)$	$(1, 0)$	0	+1	0	0	$0, +1$
Scalars								
$\chi = \begin{pmatrix} \bar{S}_1^\dagger \\ \chi^0 \end{pmatrix}$	$(4, 1, 1/2)$	$\bar{S}_1^\dagger \ (3, 1, 2/3)$ $\chi^0 \ (1, 1, 0)$	$(3, 2/3)$ $(1, 0)$	$\begin{pmatrix} +1/3 \\ -1 \end{pmatrix}$	0	+1	$\begin{pmatrix} +1/3 \\ 0 \end{pmatrix}$	$\begin{pmatrix} 0 \\ +1 \end{pmatrix}, \begin{pmatrix} -1 \\ 0 \end{pmatrix}$
H	$(1, 2, 1/2)$	$(1, 2, 1/2)$	$H_1^+ \ (1, 1)$ $H_1^0 \ (1, 0)$	0	0	0	0	0
$\Phi = \begin{pmatrix} G & R_2 \\ \bar{R}_2^\dagger & 0 \end{pmatrix}$	$(15, 2, 1/2)$	$R_2 \ (3, 2, 7/6)$ $\bar{R}_2^\dagger \ (\bar{3}, 2, -1/6)$ $G \ (8, 2, 1/2)$ $H_2 \ (1, 2, 1/2)$	$R_2^{5/3} \ (3, 5/3)$ $R_2^{2/3} \ (3, 2/3)$ $\bar{R}_2^{-1/3\dagger} \ (\bar{3}, 1/3)$ $\bar{R}_2^{2/3\dagger} \ (\bar{3}, -2/3)$ $G^+ \ (8, 1)$ $G^0 \ (8, 0)$ $H_2^+ \ (1, 1)$ $H_2^0 \ (1, 0)$	$\begin{pmatrix} 0 & +4/3 \\ -4/3 & 0 \end{pmatrix}$	0	0	$\begin{pmatrix} 0 & +1/3 \\ -1/3 & 0 \end{pmatrix}$	$\begin{pmatrix} 0 & -1 \\ +1 & 0 \end{pmatrix}$
$+\sqrt{2}T^{15}H_2$								
Gauge Bosons								
$A_\mu = \begin{pmatrix} G_\mu & X_\mu \\ X_\mu^* & 0 \end{pmatrix}$	$(15, 1, 0)$	$G_\mu \ (8, 1, 0)$ $X_\mu \ (3, 1, 2/3)$ $B'_\mu \ (1, 1, 0)$	$(8, 0)$ $(3, 2/3)$ $(1, 0)$	$\begin{pmatrix} 0 & +4/3 \\ -4/3 & 0 \end{pmatrix}$	0	0	$\begin{pmatrix} 0 & +1/3 \\ -1/3 & 0 \end{pmatrix}$	$\begin{pmatrix} 0 & -1 \\ +1 & 0 \end{pmatrix}$
W_μ	$(1, 3, 0)$	$(1, 3, 0)$	$W_\mu^\pm \ (1, \pm 1)$ $W_\mu^3 \ (1, 0)$	0	0	0	0	0
B_μ	$(1, 1, 0)$	$(1, 1, 0)$	$(1, 0)$	0	0	0	0	0

In the broken phase, this leads to the following relations between the mass matrices for the SM fermions and the underlying Yukawa matrices:

$$U_u^\dagger \hat{M}_u V_u = \frac{v_1}{\sqrt{2}} Y_1 + \frac{v_2}{2\sqrt{6}} Y_2, \quad (10)$$

$$U_d^\dagger \hat{M}_d V_d = \frac{v_1}{\sqrt{2}} Y_3 + \frac{v_2}{2\sqrt{6}} Y_4, \quad (11)$$

$$\hat{M}_e = \frac{v_1}{\sqrt{2}} Y_3 - \frac{3v_2}{2\sqrt{6}} Y_4. \quad (12)$$

Here \hat{M} are diagonal and U, V are unitary matrices describing the relation between the gauge and mass eigenstates. We work in a basis where the lepton mass matrix is flavor-diagonal. The only constraints on U 's and V 's are that $V_{\text{CKM}} = V_u^\dagger V_d$ must be reproduced. In the current study, we shall assume that all the Yukawa matrices in (9) are symmetric in the flavor space and, hence, $U_d = V_d^*$ and $U_u = V_u^*$, besides simplicity, this is motivated by the idea that this model might eventually be embedded in a variant of the minimal $SO(10)$ framework (such as proposed in Ref. [47]). Thus, we are left with just one mixing matrix V_d which we can choose freely. As we will see, this freedom is crucial for accommodating the B

anomalies without violating other constraints from lepton flavor violating observables.

In the neutrino sector we have a 9×9 complex symmetric matrix in the basis (ν, ν^c, N) which should yield the light-neutrino PMNS matrix as well as their measured mass differences.

The scalar leptoquarks in this model (R_2, \bar{R}_2 and S_1) reside in Φ and χ . After the $SU(4)_C$ breaking, the masses of the $SU(2)$ -doublet scalars conform the sum-rule

$$m_G^2 + 2m_{H'}^2 \sin^2 \beta = \frac{3}{2}(m_{R_2}^2 + m_{\bar{R}_2}^2), \quad (13)$$

where G denotes the scalar gluons and H' stands for the heavier eigenstate of the H and H_2 mixture. This, among other things, implies that one can not have both $m_{R_2}^2$ and $m_{\bar{R}_2}^2$ significantly smaller than m_G^2 and $m_{H'}^2$. It is well known [48], however, that a light scalar leptoquark with the quantum numbers of R_2 is way more suitable for a proper explanation of $R_{K^{(*)}}$ than \bar{R}_2 ; hence, in what follows, we shall work in the setting with $m_{R_2}^2 \ll m_{\bar{R}_2}^2$.

Furthermore, the mixing among the charge-2/3 components of R_2 and \bar{R}_2 with the \bar{S}_1 field emerges only from the $SU(2)_L \otimes U(1)_Y$ breaking. Hence, the physical mass eigenstates $R'_{1,2}$ are dominated

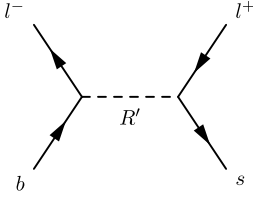


Fig. 1. Tree-level contribution to R_K via leptoquarks R' .

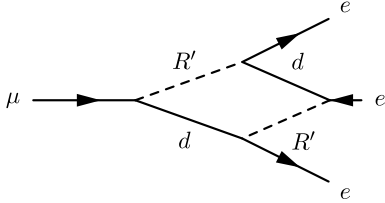


Fig. 2. One-loop contribution to $\mu \rightarrow 3e$ via leptoquarks R' .

by $R_2^{2/3}$ and $\tilde{R}_2^{2/3}$, respectively, whereas \bar{S}_1 approximates the Goldstone mode associated to the vector leptoquark.

3. Low energy observables

We turn now to a discussion of the relevant low-energy observables in the current model. First of all, we want to explain the observed deviation from the lepton universality in the B -meson decays. The Feynman diagram responsible for the tree-level contributions to R_K via the scalar leptoquarks R' is depicted in Fig. 1. It is important to notice, however, that the same leptoquark which should explain the B -meson anomalies would also contribute to other observables. At the tree level, one can expect an impact on other meson observables like $B \rightarrow X_s \nu \nu$, $B_s^0 \rightarrow \bar{l} l$ or $K_L \rightarrow e \mu$. Moreover, there are important loop contributions to $B \rightarrow X_s \gamma$, $l \rightarrow l' \gamma$ and $l \rightarrow 3l'$. An example of the responsible Feynman diagrams is shown in Fig. 2.

For our numerical study we used the `Mathematica` package `SARAH` [49–53] and extended⁴ it to support the model under consideration. In the first step, we used the model files to produce a spectrum generator based on `SPheno` [55,56]. `SPheno` calculates the mass spectrum providing the option to include all one-loop and the important two-loop corrections to neutral scalar masses [57–59] in the DR or the MS scheme. However, we are assuming here a full on-shell calculation of all masses, i.e., all shifts can be absorbed into counter-terms of the couplings leaving the mass spectrum unchanged.

In addition, `SPheno` provides an interface to `HiggsBounds` [60–62] which we used to check the constraints on the neutral scalars. Moreover, `SPheno` calculates electroweak precision as well as flavor observables. The calculation of flavor observables is based on the `FlavorKit` functionality presented in Ref. [63]. We used this feature in order to calculate the values for all necessary lepton flavor violating observables including $K_L \rightarrow e \mu$. Moreover, the values of the Wilson coefficients relevant for the B -physics calculated by `SPheno` were passed to `flavio` [64] to obtain predictions for the B -meson observables.

In Table 2 we collect the input parameters for this study. The remaining parameters affect the heavy states which do not contribute to the observables discussed below. For the fermions we take as input the known quark and lepton masses, the CKM and

Table 2

Summary of the default input values used in this analysis except if stated otherwise. All other BSM scalars have masses of the order $\mathcal{O}(m_A)$.

Numerical input values	
Y_2	$\text{diag}(10^{-8}, 10^{-7}, 10^{-5})$
Y_5	$\text{diag}(10^{-2}, 5 \cdot 10^{-2}, 10^{-1})$
$\phi_{12}, \phi_{13}, \phi_{23}$	$\pi/2, 0, \pi/4$
v_χ	$4 \cdot 10^6$ GeV
$m_A, m_{R'_1}$	$2 \cdot 10^5$ GeV, 900 GeV
$\tan\beta$	50

the PMNS matrices using the best fit values reported in [65], the Yukawa couplings Y_2 , Y_5 and $\tan\beta$. For an explanation of the B -physics observables we need an off-diagonal structure in Y_4 . We therefore take V_d as input, parametrizing it as

$$V_d = \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_{23} & s_{23} \\ 0 & -s_{23} & c_{23} \end{pmatrix} \cdot \begin{pmatrix} c_{13} & 0 & s_{13}e^{i\delta} \\ 0 & 1 & 0 \\ -s_{13}e^{i\delta} & 0 & c_{13} \end{pmatrix} \cdot \begin{pmatrix} c_{12} & s_{12} & 0 \\ -s_{12} & c_{12} & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (14)$$

with $c_{ij} = \cos\phi_{ij}$ and $s_{ij} = \sin\phi_{ij}$, and vary all three angles in the range $[0, \pi]$ with $\delta = 0$ for simplicity. In the scalar sector we take the $SU(4)_C$ -breaking VEV v_χ , the mass of the R'_1 leptoquark and the overall scale (m_A) of the heavy integer-charge Higgs bosons as input. Moreover, we fix the mass of the SM-like Higgs-boson to 125.1 GeV by adjusting λ_7 accordingly. As we take the heavy Higgs bosons to be in the multi-TeV range, we are in the decoupling limit and fulfil automatically the experimental constraints on the observed Higgs boson.

As long as R_D and R_{D^*} are concerned, there are in principle two ways to accommodate the data in the current scenario: (i) The model automatically contains a vector leptoquark of a suitable type [32]; however, its potential effects in $R_{D^{(*)}}$ are strongly suppressed by the need to be compatible with $K_L \rightarrow \mu e$ which pushes its allowed mass above 1600 TeV. (ii) A way larger effect than (i) is expected if the two scalar $SU(2)$ -doublet leptoquarks in the spectrum remain light and, at the same time, their charge $+2/3$ components entertain a large mixing. This, as discussed in Sect. 2, is impossible in the current setting. Hence, the simple model at stakes has serious issues with accommodating the existing $\bar{B} \rightarrow D^{(*)} \ell \bar{\nu}$ data.

By contrast, it is fairly easy to explain the currently observed values for R_K and R_{K^*} as demonstrated in Fig. 3. Here we have fixed the input parameters as given in Table 2 and varied one of the relevant angles. The new physics contributions to the Wilson coefficients are given by

$$C_{9,R_2}^{\text{NP}ij} = C_{10,R_2}^{\text{NP}ij} = -\frac{1}{4} \frac{1}{m_{R'_1}^2} \left(Y_4 V_d^\dagger \right)_{i3} \left(V_d Y_4^\dagger \right)_{2j}, \quad (15)$$

where $i = j = 1, 2$ for the contribution to $b \rightarrow see$ and $b \rightarrow s\mu\mu$, respectively. Using Eqs. (11) and (12) we calculate the Yukawa couplings Y_3 and Y_4 in terms of the lepton and quark masses as well as V_d . Exploiting the hierarchy in the fermion mass spectrum we get, to a good approximation,

$$C_{9,R_2}^{\text{NP}ee} \simeq -\frac{3m_b m_s}{128m_{R'_1}^2 v_{\text{ew}}^2} (1 + \tan^2\beta) f_-, \quad (16)$$

$$C_{9,R_2}^{\text{NP}\mu\mu} \simeq \frac{3m_b(m_s - m_\mu)}{128m_{R'_1}^2 v_{\text{ew}}^2} (1 + \tan^2\beta) f_+, \quad (17)$$

⁴ Details about the new feature to support unbroken subgroups in `SARAH` will be given in [54].

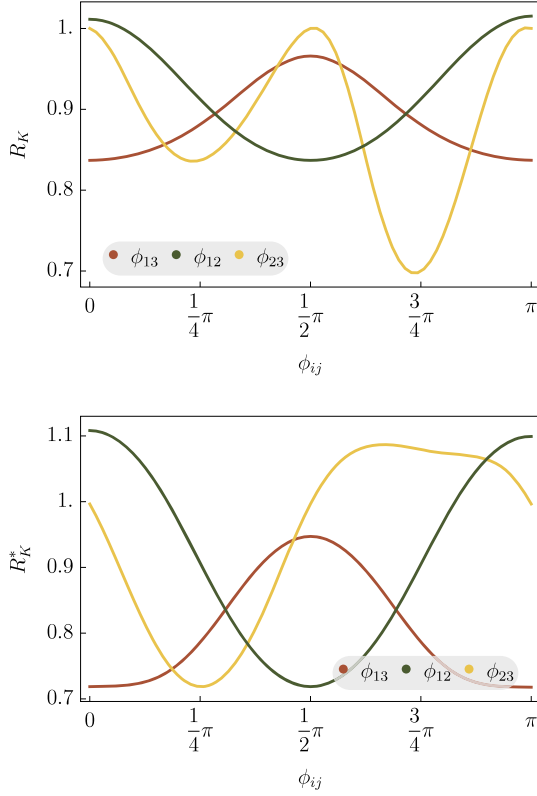


Fig. 3. R_K (first row) and R_{K^*} (second row) as a function of the mixing angles in the down-quark sector. The remaining parameters are given in Table 2.

with

$$f_{\pm} = 4 \cos 2\phi_{23} \sin 2\phi_{12} \sin \phi_{13} + \sin 2\phi_{23} \left(3 \cos 2\phi_{12} \pm 1 \pm 2 \cos 2\phi_{13} \frac{\sin^2 \phi_{12}}{\cos^2 \phi_{12}} \right). \quad (18)$$

Combining this with

$$R_K \simeq 1 + \frac{1}{|C_{9,R_2}^{\text{SM}}|^2} \left(|C_{9,R_2}^{\text{NP}\mu\mu}|^2 - |C_{9,R_2}^{\text{NP}ee}|^2 \right) \quad (19)$$

from [66] one obtains an excellent analytic approximation to the numerical results shown in Fig. 3.

In Fig. 4 we display the contour lines R_K and R_{K^*} in the $m_{R'_1}$ - $\tan\beta$ plane. The shaded regions indicate the $1\text{-}\sigma$ regions consistent with present data. The increase of $\tan\beta$ with increasing $m_{R'_1}$ can easily be understood from Eqs. (16)–(17) as the new physics contributions scale as

$$\frac{Y_4^2}{m_{R'_1}^2} \propto \frac{\tan^2 \beta}{m_{R'_1}^2}. \quad (20)$$

We recall that the leptoquarks in general also contribute to lepton-flavor violating decays of the muon such as $\mu \rightarrow e\gamma$, see for instance [67]. There are two main contributions to this observable in the current setting coming, namely, from heavy-neutrino and W -boson loops as well as leptoquark and quark loops. As an example, in Fig. 5 we show $\text{BR}(\mu \rightarrow e\gamma)$ as a function of an extra factor s rescaling the Y_2 eigenvalues in Table 2 into $Y_2 \rightarrow sY_2$. For $s \gtrsim 5$ the leptoquark loops dominate whereas for $s \leq 1$ the

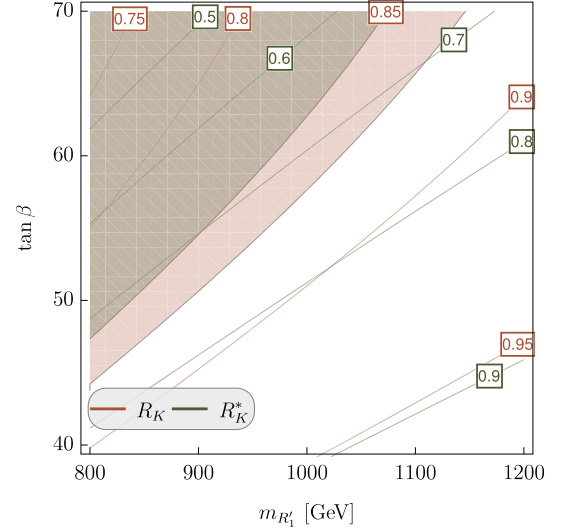


Fig. 4. Contour plots of R_K and R_{K^*} as a function of the leptoquark mass and $\tan\beta$. The shaded regions indicate the range preferred by current data. The remaining parameters are given in Table 2.

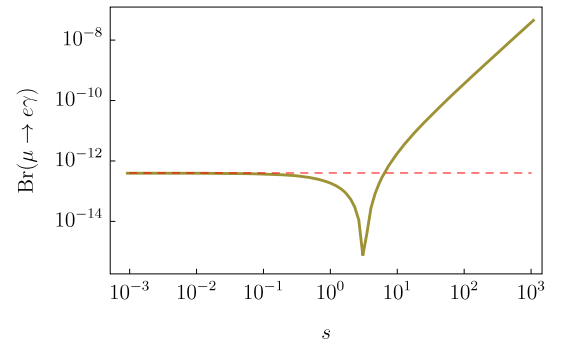


Fig. 5. $\text{BR}(\mu \rightarrow e\gamma)$ as a function of s which rescales the values of Y_2 in Table 2. The other parameters are fixed as before.

neutrino loops are dominant. The narrow minimum is due to a negative interference between both contributions.⁵

Let us point out that in the current model the bounds from $\mu \rightarrow 3e$ are in general stronger than those from $\mu \rightarrow e\gamma$ in the range interesting for R_K and R_{K^*} , see Fig. 6. The main reason for this is the negative interference in $\mu \rightarrow e\gamma$ discussed above which does neither take place in the Z -penguins nor in the box-contributions to $\mu \rightarrow 3e$ (see Fig. 2) for the same set of parameters.

In addition, we have checked that rare τ decays do not impose any constraints in the $R_{K^{(*)}}$ -interesting regions. The same holds for rare b -decays, such as $B_{s,d} \rightarrow \mu^+\mu^-$, $B \rightarrow s\gamma$, $B \rightarrow X_s\nu\nu$. Taking only the couplings of the scalar sector given in ref. [43] would yield $m_{G^{+,0}}^2 = m_{R'_1}^2 + O(v_{\text{ew}}^2)$ leading to too large contributions to ΔM_K and ΔM_B . However, this relationship gets broken by λ_{14} in Eq. (7) implying that also these bounds can be avoided.

In any case, there is another stringent constraint to be considered, namely the bound on $K_L \rightarrow e\mu$. This mode is usually used as a limit on the mass of vector leptoquarks but it is typically not being taken into account for the scalar ones. As can be seen in Fig. 7,

⁵ In principle, s (and, hence, Y_2) may be even larger than that indicated in Fig. 5 if off-diagonal elements of Y_5 were invoked together with this negative interference. However, as Y_2 does not enter the calculation of $R_{K^{(*)}}$ at the lowest order, we do not investigate this further.

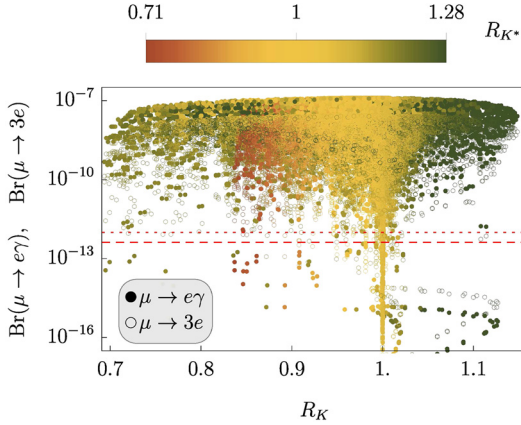


Fig. 6. Correlation between $\text{BR}(\mu \rightarrow e\gamma)$, $\text{BR}(\mu \rightarrow 3e)$ and R_K obtained by varying $\phi_{12}, \phi_{13}, \phi_{23}$ in equal steps of $\pi/27$ in the range $[0, \pi]$ each. The other parameters are fixed as given in Table 2. The color code indicates the value of R_{K^*} .

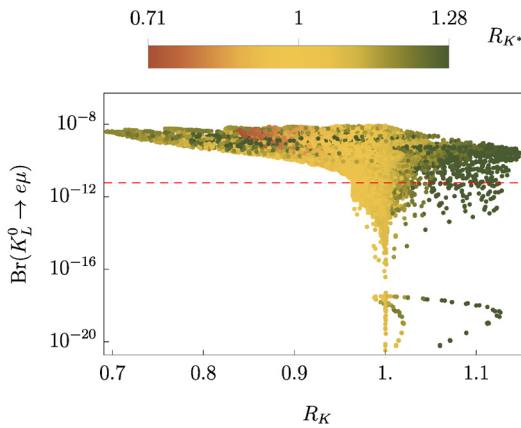


Fig. 7. Correlation between $\text{BR}(K_L \rightarrow e\mu)$ and R_K obtained by varying ϕ_{ij} as in Fig. 6. The other parameters are fixed as given in Table 2. The color code indicates the value of R_{K^*} for each point.

in the interesting region for R_K the bound is violated by several orders of magnitude, thus ruling out this model even in those tuned parts of the parameter space where all other constraints can be satisfied.

It has been argued in [32] that additional fermions in vector-like representations of the gauge group can reduce the couplings of the SM-fermions to vector leptoquarks which, in turn, may be used for lowering the generic experimental limits for their masses. The same mechanism can in principle work also for the couplings of the scalar leptoquarks such that the constraints due to $K_L \rightarrow e\mu$ can be satisfied. However, a detailed exploration of this aspect is beyond the scope of this paper and will be elaborated on elsewhere [54].

We note for completeness that leptoquarks with masses of about 1 TeV are already constrained by the LHC searches. These, however, typically focus on the situation when the decays are dominated by one channel; in the scenarios where the LQs interact through multiple couplings the corresponding bounds must be re-evaluated.

We would also like to point out that the “LQ beta-decay modes”, i.e., the decays of a heavier LQ into its lighter $SU(2)_L$ companion and W have not been considered so far within the collider searches. Nevertheless, depending on the exact mass splitting, they

may be of significant interest, especially if the on-shell W production is kinematically allowed.⁶

4. Discussion and conclusions

Motivated by the successful attempts to explain the observed values of $R_{D^{(*)}}$ and $R_{K^{(*)}}$ by leptoquarks we study a unified $SU(4)_C \otimes SU(2)_L \otimes U(1)_R$ gauge model in order to demonstrate the challenges one faces when all additional relations inherent to a unified scenario are taken into account. Among these, the dominant role is typically played by the constraints on the Yukawa couplings from the quark and lepton masses and mixing data and/or the tight connection between the relevant gauge coupling and α_s .

The model under consideration contains three different types of leptoquarks: a hypercharge-2/3 vector leptoquark and a pair of scalar leptoquarks with hypercharges 1/6 and 7/6. In its minimal version, with the SM fermion sector extended such that it supports the inverse-seesaw mechanism for neutrinos, one finds [43] that the kaon physics constrains the mass of the vector leptoquark to such an extent that it cannot significantly impact the B -physics observables.

We have shown that in the setting under consideration one can get the hypercharge-7/6 scalar leptoquark in the TeV range while, at the same time, have automatically the hypercharge-1/6 scalar leptoquark rather heavy. With this scalar sector, the low-energy effective operator structure of the model at the scale of B -mesons is such that it can accommodate R_K and R_{K^*} but neither R_D nor R_{D^*} .

We find that the allowed parameter space for R_K and R_{K^*} gets severely constrained by the bounds on rare muon decays; in particular, $\mu \rightarrow 3e$ is more important than $\mu \rightarrow e\gamma$. Neither lepton flavor violating τ -decays nor other B -physics observables lead to additional constraints. However, it turns out that no points in the available parameter space are compatible with $K_L \rightarrow e\mu$.

On the other hand, this does not imply that this kind of a leptoquark model is ruled out straight away as an explanation of $R_{K^{(*)}}$ because one can always enlarge the fermion sector by vector-like representations. In this way one may in principle reduce the couplings to the muon by mixing effects and, thus, avoid the bound due to $K_L \rightarrow e\mu$. This goes beyond the scope of this letter and will be elaborated on in a future study.

Acknowledgements

We thank H. Kolečová for discussion in the initial phase of this project. T.F., P.M. and W.P. have been supported by the DFG, project nr. PO-1337/7-1. F.S. is supported by the ERC Recognition Award ERC-RA-0008 of the Helmholtz Association. M.H. and M.M. acknowledge the support from the Grant Agency of the Czech Republic, Project No. 17-04902S.

Appendix A. Baryon number conservation

It has been noted in [43] that the model has an approximate extra $U(1)_F$ symmetry corresponding to the fermion number F , which is explicitly broken by the Majorana mass term for N . However, there is another independent accidental global symmetry $U(1)_M$, the charges of which are

⁶ Let us note that in the model of our interest the LQ-doublet mass splitting would be below m_W if we were to consider only the potential given in [43] but can be larger once the additional terms as in Eq. (7) are included.

$$\begin{aligned}
M(F_L^\alpha) &= M(\chi^\alpha) = +1, \\
M(f_{u\alpha}^c) &= M(f_{d\alpha}^c) = -1, \\
M(\Phi_\beta^\alpha) &= M(H) = M(N) = 0, \\
M(A_{\mu\beta}^\alpha) &= M(W_\mu) = M(B_\mu) = 0.
\end{aligned}
\tag{A.1}$$

Here α, β denote $SU(4)_C$ vector-like indices. Notice that we can obtain the M -charges of all the field multiplets by the prescription

$$M = (\# \text{ upper } SU(4)_C \text{ indices}) - (\# \text{ lower } SU(4)_C \text{ indices}). \tag{A.2}$$

The $U(1)_M$ symmetry of each term in the Lagrangian is then guaranteed by the fact that every upper $SU(4)_C$ index is contracted to a lower one, all carried by the dynamical fields. It is also clear that any hypothetical M -violating but $SU(4)_C$ -preserving term necessarily contains the antisymmetric tensor $\varepsilon_{\alpha\beta\gamma\delta}$. Hence, such type of a symmetry is realized in any $SU(4)_C$ model whose field content does not allow for the $SU(4)_C$ Levi-Civita symbol to occur in the interaction Lagrangian at the renormalizable level; for example, in [68], the corresponding number is called B' .

Having the M -charge at hand, we can combine it with the gauge charge (6) as

$$B = \frac{1}{4} (M + [B - L]), \tag{A.3}$$

which obviously yields the baryon number (see Table 1).

As one can verify readily, in the model under consideration both $[B - L]$ and M are spontaneously broken by $\langle \chi \rangle$ whilst their sum (A.3) remains a good symmetry even in the asymmetric phase.

On a more general ground, one can rephrase the same argument as follows: If M defined as (A.2) is a good symmetry of the unbroken-phase theory, there is no $SU(4)_C$ -Levi-Civita tensor in its Lagrangian \mathcal{L}_4 . This means that there is no $SU(3)_C$ -Levi-Civita tensor in the broken-phase Lagrangian \mathcal{L}_3 either. Consequently, a global charge defined as

$$\begin{aligned}
3B &= (\# \text{ upper } SU(3)_C \text{ indices}) \\
&\quad - (\# \text{ lower } SU(3)_C \text{ indices})
\end{aligned}
\tag{A.4}$$

generates a good symmetry of \mathcal{L}_3 and, hence, B – the usual SM baryon number – is perturbatively conserved.

Let us also note that there are two slightly different candidates for the lepton number, none of which is, however, related to a fully conserved quantity in our model. The first option is intuitive,

$$L = B - [B - L]. \tag{A.5}$$

Here, $U(1)_L$ is a good symmetry of the classical action but it is spontaneously broken, together with $[B - L]$, by $\langle \chi \rangle$. The alternative,

$$L' = F - 3B, \tag{A.6}$$

is, on the other hand, preserved by the vacuum but explicitly broken by the Majorana mass term because F is so.

References

- [1] BaBar Collaboration, J.P. Lees, et al., Phys. Rev. Lett. 109 (2012) 101802, arXiv:1205.5442.
- [2] BaBar Collaboration, J.P. Lees, et al., Phys. Rev. D 88 (2013) 072012, arXiv:1303.0571.
- [3] Belle Collaboration, A. Matyjka, et al., Phys. Rev. Lett. 99 (2007) 191807, arXiv:0706.4429.
- [4] Belle Collaboration, A. Bozek, et al., Phys. Rev. D 82 (2010) 072005, arXiv:1005.2302.
- [5] Belle Collaboration, M. Huschle, et al., Phys. Rev. D 92 (2015) 072014, arXiv:1507.03233.
- [6] LHCb Collaboration, R. Aaij, et al., Phys. Rev. Lett. 115 (2015) 111803, arXiv:1506.08614, Phys. Rev. Lett. 115 (15) (2015) 159901, Erratum.
- [7] LHCb Collaboration, R. Aaij, et al., Phys. Rev. Lett. 113 (2014) 151601, arXiv:1406.6482.
- [8] LHCb Collaboration, R. Aaij, et al., J. High Energy Phys. 08 (2017) 055, arXiv:1705.05802.
- [9] LHCb Collaboration, R. Aaij, et al., Phys. Rev. Lett. 111 (2013) 191801, arXiv:1308.1707.
- [10] F. Beaujean, C. Bobeth, D. van Dyk, Eur. Phys. J. C 74 (2014) 2897, arXiv:1310.2478, Eur. Phys. J. C 74 (2014) 3179, Erratum.
- [11] J. Lyon, R. Zwicky, arXiv:1406.0566, 2014.
- [12] W. Altmannshofer, C. Niehoff, P. Stangl, D.M. Straub, Eur. Phys. J. C 77 (2017) 377, arXiv:1703.09189.
- [13] B. Capdevila, A. Crivellin, S. Descotes-Genon, J. Matias, J. Virto, J. High Energy Phys. 01 (2018) 093, arXiv:1704.05340.
- [14] M. Bauer, M. Neubert, Phys. Rev. Lett. 116 (2016) 141802, arXiv:1511.01900.
- [15] M. Bauer, M. Neubert, Phys. Rev. D 93 (2016) 115030, arXiv:1512.06828.
- [16] A. Crivellin, D. Müller, T. Ota, J. High Energy Phys. 09 (2017) 040, arXiv:1703.09226.
- [17] W. Chao, Nucl. Phys. B 911 (2016) 231, arXiv:1512.08484.
- [18] C.W. Murphy, Phys. Lett. B 757 (2016) 192, arXiv:1512.06976.
- [19] R. Alonso, B. Grinstein, J. Martin Camalich, J. High Energy Phys. 10 (2015) 184, arXiv:1505.05164.
- [20] L. Calibbi, A. Crivellin, T. Ota, Phys. Rev. Lett. 115 (2015) 181801, arXiv:1506.02661.
- [21] S. Fajfer, N. Košnik, Phys. Lett. B 755 (2016) 270, arXiv:1511.06024.
- [22] G. Hiller, D. Loose, K. Schönwald, J. High Energy Phys. 12 (2016) 027, arXiv:1609.08895.
- [23] B. Bhattacharya, A. Datta, J.-P. Guévin, D. London, R. Watanabe, J. High Energy Phys. 01 (2017) 015, arXiv:1609.09078.
- [24] D. Buttazzo, A. Greljo, G. Isidori, D. Marzocca, J. High Energy Phys. 11 (2017) 044, arXiv:1706.07808.
- [25] J. Kumar, D. London, R. Watanabe, arXiv:1806.07403, 2018.
- [26] F.F. Deppisch, S. Kulkarni, H. Päs, E. Schumacher, Phys. Rev. D 94 (2016) 013003, arXiv:1603.07672.
- [27] R. Barbieri, A. Tesi, Eur. Phys. J. C 78 (2018) 193, arXiv:1712.06844.
- [28] R. Barbieri, G. Isidori, A. Pattori, F. Senia, Eur. Phys. J. C 76 (2016) 67, arXiv:1512.01560.
- [29] R. Barbieri, C.W. Murphy, F. Senia, Eur. Phys. J. C 77 (2017) 8, arXiv:1611.04930.
- [30] N. Assad, B. Fornal, B. Grinstein, Phys. Lett. B 777 (2018) 324, arXiv:1708.06350.
- [31] L. Di Luzio, A. Greljo, M. Nardecchia, arXiv:1708.08450, 2017.
- [32] L. Calibbi, A. Crivellin, T. Li, arXiv:1709.00692, 2017.
- [33] M. Bordone, C. Cornella, J. Fuentes-Martin, G. Isidori, Phys. Lett. B 779 (2018) 317, arXiv:1712.01368.
- [34] I. Doršner, S. Fajfer, D.A. Faroughy, N. Košnik, arXiv:1706.07779, 2017.
- [35] M. Blanke, A. Crivellin, Phys. Rev. Lett. 121 (2018) 011801, arXiv:1801.07256.
- [36] A. Greljo, B.A. Stefanek, Phys. Lett. B 782 (2018) 131, arXiv:1802.04274.
- [37] M. Bordone, C. Cornella, J. Fuentes-Martin, G. Isidori, arXiv:1805.09328, 2018.
- [38] S. Matsuzaki, K. Nishiwaki, K. Yamamoto, arXiv:1806.02312, 2018.
- [39] J.C. Pati, A. Salam, Phys. Rev. D 10 (1974) 275.
- [40] W.H. Zurek, Nature 317 (1985) 505.
- [41] H. Murayama, J. Shu, Phys. Lett. B 686 (2010) 162, arXiv:0905.1720.
- [42] A.D. Smirnov, Phys. Lett. B 346 (1995) 297, arXiv:hep-ph/9503239 [166(1995)].
- [43] P. Fileviez Perez, M.B. Wise, Phys. Rev. D 88 (2013) 057703, arXiv:1307.6213.
- [44] R.N. Mohapatra, J.W.F. Valle, Phys. Rev. D 34 (1986) 1642.
- [45] G. Valencia, S. Willenbrock, Phys. Rev. D 50 (1994) 6843, arXiv:hep-ph/9409201.
- [46] A.D. Smirnov, Mod. Phys. Lett. A 33 (2018) 1850019, arXiv:1801.02895.
- [47] S. Bertolini, L. Di Luzio, M. Malinsky, Phys. Rev. D 85 (2012) 095014, arXiv:1202.0807 [hep-ph].
- [48] G. D'Amico, et al., J. High Energy Phys. 09 (2017) 010, arXiv:1704.05438.
- [49] F. Staub, arXiv:0806.0538, 2008.
- [50] F. Staub, Comput. Phys. Commun. 181 (2010) 1077, arXiv:0909.2863.
- [51] F. Staub, Comput. Phys. Commun. 182 (2011) 808, arXiv:1002.0840.
- [52] F. Staub, Comput. Phys. Commun. 184 (2013) 1792, arXiv:1207.0906.
- [53] F. Staub, Comput. Phys. Commun. 185 (2014) 1773, arXiv:1309.7223.
- [54] T. Faber, et al., in preparation.
- [55] W. Porod, Comput. Phys. Commun. 153 (2003) 275, arXiv:hep-ph/0301101.
- [56] W. Porod, F. Staub, arXiv:1104.1573, 2011.
- [57] M.D. Goodsell, K. Nickel, F. Staub, Eur. Phys. J. C 75 (2015) 32, arXiv:1411.0675.
- [58] M. Goodsell, K. Nickel, F. Staub, Eur. Phys. J. C 75 (2015) 290, arXiv:1503.03098.
- [59] J. Braathen, M.D. Goodsell, F. Staub, arXiv:1706.05372, 2017.
- [60] P. Bechtle, O. Brein, S. Heinemeyer, G. Weiglein, K.E. Williams, Comput. Phys. Commun. 181 (2010) 138, arXiv:0811.4169.
- [61] P. Bechtle, O. Brein, S. Heinemeyer, G. Weiglein, K.E. Williams, Comput. Phys. Commun. 182 (2011) 2605, arXiv:1102.1898.
- [62] P. Bechtle, et al., Eur. Phys. J. C 74 (2014) 2693, arXiv:1311.0055.

- [63] W. Porod, F. Staub, A. Vicente, *Eur. Phys. J. C* 74 (2014) 2992, arXiv:1405.1434.
- [64] D. Straub, et al., *flav-io/flavio v0.23*, 2017.
- [65] Particle Data Group, M. Tanabashi, et al., *Phys. Rev. D* 98 (2018) 030001.
- [66] G. Hiller, M. Schmaltz, *J. High Energy Phys.* 02 (2015) 055, arXiv:1411.4773.
- [67] A. Crivellin, D. Mueller, A. Signer, Y. Ulrich, arXiv:1706.08511, 2017.
- [68] R. Foot, *Phys. Lett. B* 420 (1998) 333, arXiv:hep-ph/9708205.