Common explanation to the $R_{K^{(*)}}$, $R_{D^{(*)}}$, and ϵ'/ϵ anomalies in a 3HDM with right-handed neutrinos and connections to neutrino physics

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Scalar theories can account for the current $R_{D^{(*)}}$ measurements through a vector operator $\bar{c}_L \gamma_\mu b_L \bar{\tau}_L \gamma^\mu \nu_L$ induced at the loop level. Once the vector contribution is considered on top of a subdominant tree-level scalar component, the predicted value of $R_{D^{(*)}}$ falls within the 1σ region indicated by the experiments. We explicitly demonstrate this claim in the framework of a three Higgs doublet model extended with GeV-scale right-handed neutrinos by matching the anomalous signal for perturbative values of the involved couplings and respecting the bounds from complementary flavor physics measurements. Remarkably, we furthermore show that the proposed framework can be employed to also simultaneously explain the present $R_{K^{(*)}}$ measurement, as well as the deviation in ϵ'/ϵ currently being debated in the literature. These results are obtained by considering the contribution of relatively light right-handed neutrinos which are fundamental in mediating the processes behind the anomalous signals. In this way, our findings reveal a new possible connection that links the flavor anomalies to the phenomenology of extended Higgs sector and neutrino physics.

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I. INTRODUCTION

In recent years, B factories and the LHCb experiment have reported several anomalies [1–11] that find no satisfying explanation in the flavor structure and interactions supported by the standard model (SM) [12–40]. The interpretation of these signals in terms of new physics faces several difficulties as well; in particular, the simultaneous explanation of the $R_{K^{(*)}}$ and $R_{D^{(*)}}$ measurements proves undoubtedly challenging for the implications on further precision observables that rule out the underlying proposals [41–51].

These difficulties are grounded in the different origin of these anomalies, which seemingly require new contributions to appear at several scales and put into discussion properties, such as the universality of gauge couplings, thoroughly verified in collider experiments. For instance, since $R_{D^{(*)}}$ quantifies the ratio of processes that receive tree-level contributions within the SM, the measured value of this parameter naturally calls for new degrees of freedom below the electroweak scale. On the contrary, as $R_{K^{(*)}}$ is

purely determined by the loop structure of the SM, we naively expect that deviations in this observable originate at a much larger scale barring the fine-tuning of the involved couplings. Furthermore, the anomalous signals require new sources of lepton flavor universality violation which, in concrete frameworks, are generally tightly constrained by collider and complementary flavor experiments.

If we analyze the SM case, we see that the only source of flavor nonuniversality is the Higgs sector, where the Yukawa couplings of fermions necessarily reflect the measured mass hierarchy. Consequently, extending the Higgs sector could be a natural way to implement the additional sources of flavor nonuniversality at the basis of the measured signals.

In the present paper we pursue this possibility, showing that the $R_{K^{(*)}}$ and $R_{D^{(*)}}$ measurements can simultaneously be addressed with an extended Higgs sector. The result is based on the preliminary investigation in Ref. [31], where we demonstrated how scalar extensions of the SM can generate a sizeable vector operator $\bar{c}_L \gamma_\mu b_L \bar{\tau}_L \gamma^\mu \nu_L$ at the loop level. Explaining the mentioned flavor anomalies with respect to the complementary experimental bounds requires at least two new scalar particles, which must belong to different gauge multiplets, as well as the inclusion of right-handed neutrinos (RHNs). A natural ultraviolet completion of the proposed framework is then provided by a three Higgs doublet model (3HDM) augmented with RHNs, which we analyze in detail in the following.

As we will show, the inclusion of the tree-level scalar operators contributions allows the model to reproduce the anomalous $R_{D^{(*)}}$ signal for values of the involved couplings

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of order unity. We consider all the emerging effective chiral operators, including the ones mediated by relatively light right-handed neutrinos usually omitted in these studies. We explicitly propose the Yukawa coupling texture of the additional Higgs doublets needed to explain the anomaly. Remarkably, we find that the same texture also allows for the explanation of the $R_{K^{(*)}}$ anomaly owing to the presence of Majorana RHNs used here exclusively to mediate the processes behind these signals.¹ In this light, flavor anomalies could therefore constitute the first indirect collider signature supporting the existence of RHNs. The proposed framework then connects the issue of flavor measurements to other open questions within neutrino physics and cosmology, such as the neutrino mass generation mechanism and the dynamics of baryogenesis via leptogenesis, extending the phenomenological reach of flavor experiments. We highlight that RHNs in the considered mass range can be discovered by the planned SHIP experiment at CERN [57], as well as in collider experiments through the implied lepton number violating effects [58].

A brief part of the present paper is dedicated to the assessment of a further potential flavor anomaly connected to the amount of direct CP violation measured in $K \rightarrow \pi\pi$ decays and quantified in the ratio ϵ'/ϵ . In this case, without advocating for the necessity of a new physics contribution, we find that a minimal modification of the proposed Yukawa texture would allow us to explain the signal should it survive further scrutiny by the dedicated community.

The paper is organized as follows. Section II is dedicated to the $R_{D^{(*)}}$ anomaly: After a brief review of its experimental status, we show how the considered framework explains the anomalous signal. The $R_{K^{(*)}}$ anomaly is studied in a similar fashion in Sec. III. In Sec. IV we then present the result of our investigation, analyze the compatibility of the two anomalies, and detail the potential impact of new determinations of ϵ'/ϵ . Finally, our conclusions are presented in Sec. V.

II. THE $R_{D^{(*)}}$ ANOMALY

We start by briefly reviewing the present experimental status of the $R_{D^{(*)}}$ anomaly, setting out the formalism at the basis of the analysis.

A. Experimental status and effective Lagrangian

The results of the LHCb experiment [1,2] have underlined the presence of lepton flavor violating dynamics in the charged decays of the B meson. In more detail, anomalies have been measured in the ratio of branching fractions

¹ Alternative explanations of the anomalous signals which rely on the presence of RHNs in the involved final states can be found in Refs. [27,28,52–56].

$$R_{D^{(*)}} = \frac{\mathcal{B}(\bar{B} \rightarrow D^{(*)} \tau \bar{\nu})}{\mathcal{B}(\bar{B} \rightarrow D^{(*)} \ell \bar{\nu})} \quad (1)$$

for $\ell = e, \mu$. Being defined as ratios, R_D and R_{D^*} are observables of particular importance that test lepton universality in the transition $b \rightarrow c \tau \bar{\nu}_\tau$ net of the involved hadronic uncertainties.

If considered along the initial measurements by the BABAR [3,4] and Belle [5–7] Collaborations, the LHCb observations set [59]

$$\begin{aligned} R_D^{\text{exp}} &= 0.407 \pm 0.039 \pm 0.024, \\ R_{D^*}^{\text{exp}} &= 0.306 \pm 0.013 \pm 0.007 \end{aligned} \quad (2)$$

highlighting a strong affinity for third-generation leptons. The implied departure from universality falls well beyond the SM predictions

$$\begin{aligned} R_D^{\text{SM}} &= 0.300 \pm 0.008, \\ R_{D^*}^{\text{SM}} &= 0.252 \pm 0.003, \end{aligned} \quad (3)$$

with a combined significance of about 4σ [59].

The possible presence of new physics in B anomalies is investigated by detailing the low-energy regime of a potential model through the construction of an effective theory characterized by dimension-six operators that preserve color and electric charge. Then, at the b -quark mass scale, the theory is described by the effective Lagrangian [60–63]

$$\begin{aligned} \mathcal{L}_{\text{eff}}^{b \rightarrow c \ell \bar{\nu}} &= -\frac{2G_F V_{cb}}{\sqrt{2}} \left(C_{VL}^\ell \mathcal{O}_{VL}^\ell + C_{AL}^\ell \mathcal{O}_{AL}^\ell \right. \\ &\quad \left. + C_{SL}^\ell \mathcal{O}_{SL}^\ell + C_{PL}^\ell \mathcal{O}_{PL}^\ell \right), \end{aligned} \quad (4)$$

where $\ell = e, \mu, \tau$. The four effective operators that appear above are given by

$$\mathcal{O}_{VL}^\ell = [\bar{c} \gamma^\mu b] [\bar{\ell} \gamma_\mu P_L \nu_\ell], \quad (5)$$

$$\mathcal{O}_{AL}^\ell = [\bar{c} \gamma^\mu \gamma_5 b] [\bar{\ell} \gamma_\mu P_L \nu_\ell], \quad (6)$$

$$\mathcal{O}_{SL}^\ell = [\bar{c} b] [\bar{\ell} P_L \nu_\ell], \quad (7)$$

$$\mathcal{O}_{PL}^\ell = [\bar{c} \gamma_5 b] [\bar{\ell} P_L \nu_\ell]. \quad (8)$$

In this formalism, the dynamics of the $b \rightarrow c \ell \bar{\nu}$ process is fully encoded in a corresponding set of Wilson coefficients $C^\ell = C_{VL}^\ell, C_{AL}^\ell, C_{SL}^\ell, C_{PL}^\ell$, which parametrize the relevant observables through the involved branching fractions

$$R_{D^{(*)}} = \frac{\mathcal{B}_{\tau}^{D^{(*)}}(m_\tau, \mathbf{C}^\tau)}{\mathcal{B}_{\mu/e}^{D^{(*)}}(m_{\mu/e}, \mathbf{C}^{\mu/e})}. \quad (9)$$

The contributions of heavy states to the effective Lagrangian in Eq. (4) is computed by matching the full theory to the effective one at the scale where the heavy degrees of freedom are integrated out. The final expression of the Lagrangian is then obtained solving the renormalization group equations down to the scale at which the process is probed. The matching has been performed diagrammatically, and we checked that the QCD and QED corrections induced by the renormalization group equations evolution have a negligible impact on the couplings in Eq. (13).

A general analysis of Wilson coefficients provides the first portrait of the framework we seek. For instance, the simplest scalar extensions of the SM yield tree-level contributions to the scalar and pseudoscalar Wilson coefficients C_{SL}^r and C_{PL}^r which, in principle, allow us to explain the anomalous B -physics signal. As shown in the left panel of Fig. 1, new contributions to these quantities let the theoretical predictions enter the 68% confidence interval associated to the signal denoted by the red areas in the plot. Such a solution, however, is invalidated by measurements of the B_c lifetime, which severely constrain the pseudoscalar Wilson coefficient due to the mass hierarchy of the SM:

$$\mathcal{B}_{\tau\nu} = \frac{m_{B_c} m_\tau^2 f_{B_c}^2 G_F^2 |V_{cb}|^2}{8\pi\Gamma_{B_c}} \left(1 - \frac{m_\tau^2}{m_{B_c}^2}\right)^2 \times \left| \frac{m_{B_c}^2}{m_\tau(m_b + m_c)} C_{PL}^r - C_{AL}^r \right|^2. \quad (10)$$

The impact of this constraint is represented in both panels of Fig. 1 by the areas shaded in light and dark gray, which

indicate the values of the Wilson coefficients that result in deviations larger than 10% or 30% from the measured B_c lifetime, respectively. As we can see, solutions characterized by large values of C_{PL}^r are obviously disfavored.

As originally proposed in Ref. [31], a possible way for scalar extensions to handle the B_c lifetime constraint is to rely on scalar loop contributions to the vector and pseudovector Wilson coefficients. The case is shown in the right panel of Fig. 1, which makes clear that modest values of C_{VL}^r and C_{AL}^r allow us to explain the observed anomaly even when the strongest constraint on the B_c lifetime is considered. A potential issue with this solution is that contributions to C_{VL}^r and C_{AL}^r are generated in scalar theories only at the loop level. The required magnitude of about $\mathcal{O}(10^{-1})$ may then impose couplings on the verge of nonperturbativity [31].

We address this problem here by explaining the observed violation of lepton universality through vector and pseudovector contributions on top of a subdominant pseudoscalar component used to relax the values of the coupling involved in the loop diagrams that generate C_{VL}^r and C_{AL}^r . The interplay between the involved Wilson coefficients is shown in the right panel of Fig. 1, where crescent values of C_{PL}^r allow an excellent fit of the anomaly for lower values of the vector and pseudovector contributions.

B. The 3HDM contribution to $R_{D^{(*)}}$

In order to respect the tight bounds on $B \rightarrow X_s \gamma$, as well as further constraints imposed by penguin diagrams, we follow the setup of Ref. [31] and consider three different scalar $SU(2)$ doublets

FIG. 1. Model-independent fit of the anomalous signal as a function of the indicated Wilson coefficients from Eq. (4). The 68%, 95%, and 99% confidence intervals selected by the joint fit of R_D and R_{D^*} are marked by the red, orange, and yellow areas, respectively. The shaded light (dark) gray areas show instead the current bound from the measured B_c lifetime assuming a 10% (30%) maximal allowed deviation.

FIG. 2. Additional diagrams for the process $b \rightarrow c\tau\bar{\nu}$ supported by the considered 3HDM. The first diagram sources the tree-level expression for C_{SL}^τ and C_{PL}^τ given in Eqs. (14) and (15). The last two diagrams, instead, yield the contributions to C_{VL}^τ and C_{AL}^τ reported in Eqs. (16) and (17). We denoted with the symbol ν an active SM neutrino.

$$H_0 = \begin{pmatrix} H_0^+ \\ H_0^0 \end{pmatrix}, \quad H_1 = \begin{pmatrix} H_1^+ \\ H_1^0 \end{pmatrix}, \quad H_2 = \begin{pmatrix} H_2^+ \\ H_2^0 \end{pmatrix}, \quad (11)$$

with H_0 being the SM Higgs doublet. We arrange the scalar potential so that the new doublets do not develop vacuum expectation values, and restrict their masses to the $\sim 300\text{--}350$ GeV range to remain within the reach of current collider searches. As mentioned before, we also consider three RH neutrinos ν_R with corresponding Majorana mass terms. In this regard, we require that $m_b - m_c - m_\tau \lesssim m_{\nu_R}$ to prevent the related b decay channel, maintaining however $m_{\nu_R} \ll m_t$ to retain sizeable loop contributions. From the neutrino physics point of view, these additional states

are therefore sufficiently heavy to decouple from the low-energy dynamics of active SM neutrinos.

The interactions of the new scalar states are detailed in the following Lagrangian

$$\begin{aligned} -\mathcal{L} \supset & \bar{Q}_L \tilde{H}_1 \mathcal{Y}_1^u u_R + \bar{Q}_L \tilde{H}_2 \mathcal{Y}_2^u u_R + \bar{L}_L H_1 \mathcal{Y}_1^\ell \ell_R \\ & + \bar{L}_L H_2 \mathcal{Y}_2^\ell \ell_R + \bar{L}_L \tilde{H}_1 \mathcal{Y}_1^\nu \nu_R + \bar{L}_L \tilde{H}_2 \mathcal{Y}_2^\nu \nu_R \\ & + \bar{Q}_L H_1 \mathcal{Y}_1^d d_R + \text{H.c.}, \end{aligned} \quad (12)$$

where, for sake of minimality, we set $\mathcal{Y}_2^d = 0$.

The Yukawa texture that we investigate is of the form²

$$\begin{aligned} \mathcal{Y}_1^u &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & f_{\tilde{c}_L t_R} \\ 0 & f_{\tilde{b}_L c_R} & 0 \end{pmatrix}, & \mathcal{Y}_1^d &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & f_{\tilde{c}_L b_R} \\ 0 & 0 & 0 \end{pmatrix}, & \mathcal{Y}_2^u &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & g_{\tilde{b}_L t_R} \end{pmatrix}, \\ \mathcal{Y}_1^\nu &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & f_{\tilde{\nu}_L \nu_R} & 0 \\ 0 & 0 & f_{\tilde{\nu}_L \nu_R}^t \end{pmatrix}, & \mathcal{Y}_1^\ell &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & f_{\tilde{\tau}_L \tau_R} \end{pmatrix}, & \mathcal{Y}_2^\nu &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & g_{\tilde{\nu}_L \nu_R} & 0 \\ 0 & 0 & g_{\tilde{\nu}_L \nu_R}' \end{pmatrix}, & \mathcal{Y}_2^\ell &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & g_{\tilde{\nu}_L \tau_R} \end{pmatrix}. \end{aligned} \quad (13)$$

In the above equations, as well as throughout the rest of the paper, the symbol f denotes the coupling of the H_1 doublet with the fields indicated by the subscript. Analogously, we indicate with g the couplings of H_2 . The elements of the new Yukawa matrices rendered in violet regulate the contribution to the pseudoscalar Wilson coefficient sourced by the first diagram in Fig. 2. The elements in teal enter, instead, the vector and pseudovector Wilson coefficient

²In the absence of a symmetry forbidding the vanishing elements, radiative corrections give rise to nonzero contributions that modify the proposed texture. It is possible to show that as long as the vanishing element receives contributions below about 10% of the remaining entries, our mechanism still provides a compelling explanation for the observed flavor anomalies. For instance, the constraint due to $b \rightarrow s\gamma$ shown in Fig. 4 excludes only unrealistically large radiative corrections of the order of 20% of the considered quark Yukawa couplings.

through the remaining loop diagrams. Notice that $f_{\tilde{\tau}_L \tau_R}$ is involved in both contributions, depending on the considered component of H_1 . Lastly, the terms in orange affect exclusively the computation of the $R_{K^{(*)}}$ anomaly presented in Sec. III. For the sake of simplicity, we assume real couplings and set to zero the remaining elements of the Yukawa matrices, with the understanding that their values are negligible within our effective description.

By integrating out the degrees of freedom above the b -quark mass scale and matching the Lagrangian in Eq. (12) to the effective one in Eq. (4), we identify the following tree-level contributions to the scalar and pseudoscalar Wilson coefficient:

$$C_{SL}^\tau = -\frac{2m_W^2}{V_{cb}g_W^2 m_{H_1^-}^2} f_{\tilde{\nu}_L \tau_R} (f_{\tilde{b}_L c_R} - f_{\tilde{c}_L b_R}), \quad (14)$$

$$C_{PL}^\tau = -\frac{2m_W^2}{V_{cb}g_w^2m_{H_1^-}^2}f_{\bar{\nu}_L\tau_R}(f_{\bar{b}_Lc_R} + f_{\bar{c}_Lb_R}), \quad (15)$$

where g_w is the coupling constant of SM weak interactions. Notice that gauge invariance imposes $f_{\bar{\nu}_L\tau_R} \equiv f_{\bar{\tau}_L\tau_R}$. Equations (14) and (15) make clear the choice of the considered Yukawa pattern, which allows $f_{\bar{b}_Lc_R}$ and $f_{\bar{c}_Lb_R}$ to separately source the tree-level contribution. In this way, two independent degrees of freedom regulate C_{SL} and C_{PL} , making it possible to exploit the Cabibbo-Kobayashi-Maskawa matrix enhancement to maintain perturbative values of the involved couplings.

The contributions to the vector and pseudovector Wilson coefficient due to the loop diagram in Fig. 2 amount instead to

$$\begin{aligned} C_{VL}^{\tau(1)} &= -C_{AL}^{\tau(1)} \\ &= \left(-\frac{m_W^2}{8\pi^2 V_{cb}g_w^2}\right)(f_{\bar{c}_Lt_R}g_{\bar{b}_Lt_R}f'_{\bar{\nu}_L\nu_R}g'_{\bar{\nu}_L\nu_R}) \\ &\quad \times D_{dd00}[m_{\nu_R}^2, m_t^2, m_{H_1^0}^2, m_{H_2^-}^2], \end{aligned} \quad (16)$$

$$\begin{aligned} C_{VL}^{\tau(2)} &= -C_{AL}^{\tau(2)} \\ &= \left(-\frac{m_W^2}{8\pi^2 V_{cb}g_w^2}\right)(f_{\bar{c}_Lt_R}g_{\bar{b}_Lt_R}f_{\bar{\tau}_L\tau_R}g_{\bar{\nu}_L\tau_R}) \\ &\quad \times D_{dd00}[m_t^2, m_{H_1^0}^2, m_{H_2^-}^2, m_\tau^2], \end{aligned} \quad (17)$$

where we used $g_{\bar{\tau}_L\nu_R} = g'_{\bar{\nu}_L\nu_R}$ and indicated with D_{dd00} the four-point loop integral:

$$\begin{aligned} D_{dd00}[m_1^2, m_2^2, m_3^2, m_4^2] \\ = \frac{(2\pi\mu)^{4-D}}{4i\pi^2} \int d^D q \frac{q^2}{(q^2 - m_1^2)(q^2 - m_2^2)(q^2 - m_3^2)(q^2 - m_4^2)}. \end{aligned} \quad (18)$$

Clearly, the final expressions for the Wilson coefficients are obtained as $C_{VL}^\tau = C_{VL}^{\tau(1)} + C_{VL}^{\tau(2)}$ and $C_{AL}^\tau = C_{AL}^{\tau(1)} + C_{AL}^{\tau(2)}$.

We remark that the introduction of the $f_{\bar{c}_Lb_R}$ coupling sources $b \rightarrow s$ tree-level transitions that seem to induce an unacceptably large contribution to the B_s mixing. However, the absence of CP -violating terms in the scalar potential of the new Higgs doublets imposes the mass degeneracy of the CP -odd and -even components in a way that the overall tree-level contribution to the $(\bar{s}\hat{X}b)^2$ operator actually vanishes. As for next-to-leading-order contributions, we expect the presence of a subdominant correction as large as 10%–20% of the SM contribution. According to the analysis in Ref. [64], our solution could be further investigated by tracking the impact of possible phases on the proposed Yukawa texture, which were neglected here because they were irrelevant to the anomalies under discussion.

FIG. 3. 3HDM contributions to $b \rightarrow s\gamma$. Notice that the diagrams involve one Higgs doublet at a time.

Before moving to the discussion of the $R_{K^{(*)}}$ anomaly, we address below the main experimental bounds that the proposed solution for $R_{D^{(*)}}$ faces.

C. Main experimental constraints

The main experimental bounds that oppose the proposed solution for the $R_{D^{(*)}}$ anomaly are due to measurements of $B \rightarrow X_s\gamma$, $B^+ \rightarrow K^+\nu\bar{\nu}$, and $B^0 \rightarrow K^{*0}\nu\bar{\nu}$.

The introduction of extra $SU(2)$ doublets through their Yukawa interactions affects the dimension-five photon and gluon dipole operators

$$\begin{aligned} P_7 &= \frac{e}{16\pi^2} m_b (\bar{s}\sigma^{\mu\nu}P_R b) F_{\mu\nu}, \\ P_8 &= \frac{g_3}{16\pi^2} m_b (\bar{s}\sigma^{\mu\nu}T^a P_R b) G_{\mu\nu}^a, \end{aligned} \quad (19)$$

and modifies via the effective Lagrangian [65]

$$\mathcal{L}_{\text{eff}} = \frac{4G_F}{\sqrt{2}} V_{ts}^* V_{tb} (C_7 P_7 + C_8 P_8) \quad (20)$$

the SM prediction for $B \rightarrow X_s\gamma$ [65–67].

It is then clear that measurements of this quantity limit possible new physics contributions resulting from Eq. (13). Employing comparable couplings for both extra doublets to induce an enhancement of the box contribution to $R_{D^{(*)}}$ of about a factor of 4, in particular, generates unacceptable large contributions to $B \rightarrow X_s\gamma$ from the diagrams in Fig. 3. This is manifest even in the limit where one of the two scalar doublets does not participate in the dynamics of the anomaly (2HDM limit³), resulting in the bound on the new contribution to C_{VL}^τ from $B \rightarrow X_s\gamma$ plotted in Fig. 4.

The bound therefore provides a clear indication in favor of our mechanism, which relies on separate gauge multiplets with complementary roles that generate the required amount of lepton flavor universality violation without inducing large corrections to the photon and Z vertices. We remark that, within the full model, it is in

³In this case, we implicitly extend the Yukawa couplings of the active doublet to include the interactions ascribed to the excluded field within the full 3HDM.

FIG. 4. The contribution of the model to C_{VL}^r in the 2HDM limit, where neither of the new scalar doublets participate in the dynamics of the anomaly. The green and orange points are, respectively, allowed and excluded by their present measurements of $B \rightarrow X_s \gamma$. The range of masses and couplings adopted in the scan for the active extra Higgs doublet is the same as in the full 3HDM framework.

principle possible to adjust the value of $\mathcal{Y}_{13,3}^d$ to cancel the new large corrections to $b \rightarrow s\gamma$ without suppressing, at the same time, the box diagrams controlling the anomaly. However, we do not pursue such a possibility because of the amount of fine-tuning implied.

As we have seen, measurements of $B \rightarrow X_s \gamma$ yield important constraints on the Yukawa couplings of the new Higgs doublet within the quark sector. In order to investigate similar constraints that potentially target the leptonic sector, we now turn our attention to $b \rightarrow s\nu_L \bar{\nu}_L$. Our 3HDM indeed produces a similar enhancement in the related processes via dimension-six four-fermion operators

$$O_{L/R}^i = \frac{e^2}{8\pi^2} (\bar{s} \gamma_\mu P_{L/R} b) (\bar{\nu}^i \gamma^\mu P_L \nu^i) \quad (21)$$

encoded in the effective Lagrangian

$$\mathcal{L}_{\text{eff}} = \frac{4G_F}{\sqrt{2}} V_{ts}^* V_{tb} \sum_{i=e, \mu, \tau} (C_L^i O_L^i + C_R^i O_R^i). \quad (22)$$

The main contribution is due to the diagram in Fig. 5, which leads to a new flavor violating deformation of C_L^r given by

$$C_L^r = - \left(\frac{m_W^2 V_{cs}}{e^2 g_w^2 V_{ts}^*} \right) (f_{\tilde{c}_L t_R} g_{\tilde{b}_L t_R} f_{\tilde{\tau}_L \tau_R} g_{\tilde{\nu}_L \tau_R}) \times D_{dd00}[m_{H_1^-}^2, m_{H_2^-}^2, m_t^2]. \quad (23)$$

Notice that the absence of RHNs in the final states prevents $f_{\tilde{\nu}_L \nu_R}^l$ and $g_{\tilde{\nu}_L \nu_R}^l$ from bearing effects on the dynamics of

FIG. 5. The main 3HDM contribution to $b \rightarrow s\nu\bar{\nu}$.

$b \rightarrow s\nu\bar{\nu}$. The current upper bounds on the branching ratios are

$$\begin{aligned} \mathcal{B}_K &= \text{BR}(B^+ \rightarrow K^+ \nu \bar{\nu}) < 1.7 \times 10^{-5}, \quad [68] \\ \mathcal{B}_{K^*} &= \text{BR}(B^+ \rightarrow K^{*+} \nu \bar{\nu}) < 4.0 \times 10^{-5}, \quad [69] \end{aligned} \quad (24)$$

and supported by the SM results

$$\begin{aligned} \mathcal{B}_K^{\text{SM}} &= \text{BR}(B^+ \rightarrow K^+ \nu \bar{\nu}) = (3.98 \pm 0.43 \pm 0.19) \times 10^{-6}, \\ \mathcal{B}_{K^*}^{\text{SM}} &= \text{BR}(B^+ \rightarrow K^{*+} \nu \bar{\nu}) = (9.19 \pm 0.86 \pm 0.50) \times 10^{-6} \end{aligned} \quad (25)$$

have been used in Ref. [68] to derive the following 90% C.L. upper bounds

$$\frac{\mathcal{B}_K}{\mathcal{B}_K^{\text{SM}}} < 4.3, \quad \frac{\mathcal{B}_{K^*}}{\mathcal{B}_{K^*}^{\text{SM}}} < 4.4. \quad (26)$$

We consider the impact of these constraints in the numerical analysis presented in Sec. IV.

III. THE $R_{K^{(*)}}$ ANOMALY

We now turn our attention to the $R_{K^{(*)}}$ anomaly, discussing its experimental status and showing how it can be addressed in the present framework.

A. Experimental status and effective Lagrangian

Another longstanding anomaly highlighted by B -physics experiments concerns the neutral current transition $b \rightarrow s \ell^+ \ell^-$. In more detail, the LHCb experiment has found anomalous values of

$$R_K = \frac{\mathcal{B}(B^+ \rightarrow K^+ \mu^+ \mu^-)}{\mathcal{B}(B^+ \rightarrow K^+ e^+ e^-)} \quad (27)$$

and

$$R_{K^*} = \frac{\mathcal{B}(B^0 \rightarrow K^{*0} \mu^+ \mu^-)}{\mathcal{B}(B^0 \rightarrow K^{*0} e^+ e^-)} \quad (28)$$

reporting, respectively [8,9],

$$R_K^{\text{exp}} = 0.745_{-0.074}^{+0.090} \pm 0.036 \quad \text{for } 1 \text{ GeV}^2 \leq q^2 \leq 6 \text{ GeV}^2, \quad (29)$$

and

$$R_{K^*}^{\text{exp}} = \begin{cases} 0.66_{-0.07}^{+0.11} \pm 0.03 & \text{for } 0.045 \text{ GeV}^2 \leq q^2 \leq 1.1 \text{ GeV}^2, \\ 0.69_{-0.07}^{+0.11} \pm 0.05 & \text{for } 1.1 \text{ GeV}^2 \leq q^2 \leq 6 \text{ GeV}^2, \end{cases} \quad (30)$$

where q^2 is the invariant mass of the final state dilepton system.

The corresponding SM predictions⁴ suppressed by the Glashow-Iliopoulos-Maiani mechanism amount to

$$R_K^{\text{SM}} = 1.0004 \pm 0.0002 \quad (31)$$

and

$$R_{K^*}^{\text{SM}} = \begin{cases} 0.926 \pm 0.003, \\ 0.9965 \pm 0.0005, \end{cases} \quad (32)$$

consequently giving rise to a discrepancy with a significance of about 5σ [69] depending on the details of the fit.

The low-energy effective Lagrangian describing the $b \rightarrow s \ell \ell$ transition is

$$-\mathcal{L}_{bs} = -\frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* \frac{e^2}{16\pi^2} \sum_i (C_i \mathcal{O}_i + C_i' \mathcal{O}_i'), \quad (33)$$

where the dimension-six operators $\mathcal{O}_i$ are defined as

$$\begin{aligned} \mathcal{O}_7 &= (\bar{s} P_L b)(\bar{l} l), & \mathcal{O}_7' &= (\bar{s} P_R b)(\bar{l} l), \\ \mathcal{O}_8 &= (\bar{s} P_L b)(\bar{l} \gamma_5 l), & \mathcal{O}_8' &= (\bar{s} P_R b)(\bar{l} \gamma_5 l), \\ \mathcal{O}_9 &= (\bar{s} \gamma_\mu P_L b)(\bar{l} \gamma^\mu l), & \mathcal{O}_9' &= (\bar{s} \gamma_\mu P_R b)(\bar{l} \gamma^\mu l), \\ \mathcal{O}_{10} &= (\bar{s} \gamma_\mu P_L b)(\bar{l} \gamma^\mu \gamma_5 l), & \mathcal{O}_{10}' &= (\bar{s} \gamma_\mu P_R b)(\bar{l} \gamma^\mu \gamma_5 l). \end{aligned} \quad (34)$$

Global fits of the anomalies presently converge on preferred sets of Wilson coefficients [22,69–71], with the highest pull given by $C_9 - C_9^{\text{SM}} \simeq -1.21$ or $C_9 - C_9^{\text{SM}} = -(C_{10} - C_{10}^{\text{SM}}) \simeq -0.67$ that sets an almost 5σ discrepancy with respect to the SM predictions.

⁴As computed with the FLAVIO-0.21.2 package [69].

FIG. 6. The diagram responsible for reproducing the $R_{K^{(*)}}$ anomaly within the considered 3HDM.

B. Explaining $R_{K^{(*)}}$ in the 3HDM

The explanation of the $R_{K^{(*)}}$ anomaly in our 3HDM reflects, for the most part, the steps required to account for the $R_{D^{(*)}}$ one. The possible role of the scalar operators is once again strongly constrained, in this case by observations of $B_s \rightarrow \mu^+ \mu^-$. We therefore seek an enhancement of the $V - A$ quark current, along with a suppression of the scalar Wilson coefficient in order to fit the current measurements.

The explanations of the two anomalies share a common origin in the couplings of the quark sector $f_{\tilde{c}_L t_R}$ and $g_{\tilde{b}_L t_R}$. The additional source of lepton flavor universality violation is implemented here through the interactions with the RHNs and muons

$$-\mathcal{L} \supset f_{\tilde{\mu}_L \nu_R} H_1^- \bar{\mu}_L \nu_R + g_{\tilde{\mu}_L \nu_R} H_2^- \bar{\mu}_L \nu_R + \text{H.c.}, \quad (35)$$

where $f_{\tilde{\mu}_L \nu_R} \equiv f_{\tilde{\nu}_L \nu_R}$, $g_{\tilde{\mu}_L \nu_R} \equiv g_{\tilde{\nu}_L \nu_R}$.

Our choice of couplings for the charged scalar fields results in vectorial Wilson coefficients that obey $C_9 = -C_{10}$, with the contribution of the diagram in Fig. 6 amounting to

$$\begin{aligned} -C_9 &= C_{10} \\ &= -\frac{m_W^2}{4\pi\alpha g_w^2} \frac{V_{cs}^*}{V_{ts}} (f_{\tilde{c}_L t_R} g_{\tilde{b}_L t_R} f_{\tilde{\mu}_L \nu_R} g_{\tilde{\mu}_L \nu_R}) \\ &\quad \times D_{dd00}[m_t^2, m_{H_1^-}^2, m_{H_2^-}^2], \end{aligned} \quad (36)$$

where D_{dd00} is the same four-point scalar integral given in Eq. (18). Notice that the $t_R \tilde{s}_L$ coupling is obtained from the $t_R \tilde{c}_L$ one via the Cabibbo-Kobayashi-Maskawa element V_{cs}^* .

IV. RESULTS

Here we gather the results obtained for the $R_{D^{(*)}}$ and $R_{K^{(*)}}$ anomalies through the analysis detailed in the previous sections, discuss their compatibility within the present framework, and remark on the possible impact of new determinations of ϵ'/ϵ on our conclusions.

FIG. 7. The $R_{D^{(*)}}$ anomaly: Benchmark points obtained with the considered 3HDM. The red, orange, and yellow areas, respectively, indicate the 68%, 95%, and 99% confidence intervals selected by the joint fit of R_D and R_{D^*} . The shaded light (dark) gray areas represent instead the current bound from the measured B_c lifetime assuming a 10% (30%) maximal allowed deviation. The two panels differ by the indicated value of the vector Wilson coefficients.

A. Numerical analysis

As for the anomalies, the joint effect of the tree-level scalar contribution and of the operator $[\bar{c}\gamma^\mu P_L b][\bar{\ell}\gamma_\mu P_L \nu_\ell]$ generated by the same particles at the one-loop level allows access to the 2σ and 1σ $R_{D^{(*)}}$ regions for values of the Yukawa sector well within the perturbative regime. Remarkably, once supplemented with an extra coupling to muons, we find that the same interactions between quarks and the new scalar doublets allow us to also explain the $R_{K^{(*)}}$ signal.

The result concerning the $R_{D^{(*)}}$ anomaly is illustrated in isolation in Fig. 7. Here we show four benchmark points in the space of Wilson coefficients obtained through the contributions detailed in Sec. II. The points in red comply with the 10% lifetime bounds on the B_c decay, while the blue ones refer to the corresponding 30% limit. The areas shaded in red, orange, and yellow indicate the 68%, 95%, and 99% confidence interval indicated by current measurements, respectively. The two panels differ by the considered value of the vector Wilson coefficients $C_{VL}^r = -C_{AL}^r$ induced by the 3HDM at the one-loop level. As shown in the right panel, larger values of this quantity allow us to fit the anomaly with a contribution from C_{PL}^r small enough to comply with the B_c lifetime bounds.

The benchmark points have been obtained by maximizing the one-loop contributions encoded in the vector and pseudovector operators, which depend on the set of couplings rendered in teal in Eq. (13). Considering the collider phenomenology analysis presented in Ref. [31], we set the values of the new quark Yukawa couplings to $f_{\bar{c}_L t_R} = 0.8$ and $g_{\bar{b}_L t_R} = 0.8$ or $g_{\bar{b}_L t_R} = 1$. The magnitude of

the RHN and lepton couplings $f'_{\bar{\nu}_L \nu_R} = g'_{\bar{\nu}_L \nu_R}$ and $f_{\bar{\tau}_L \tau_R} = g_{\bar{\nu}_L \tau_R}$ shown in Fig. 8 is then obtained by setting the Wilson coefficients to the indicated values.

The subdominant tree-level contribution of C_{SL}^r and C_{PL}^r is subsequently obtained through Eq. (14), in compliance with the B_c lifetime bounds. The required magnitude of the involved couplings $f_{\bar{b}_L c_R}$ and $f_{\bar{c}_L b_R}$, which do not enter the vector contribution, is presented in Fig. 9. Here we let the mass of the new scalar particles vary in the selected range and set $f_{\bar{\nu}_L \tau_R} \equiv f_{\bar{\tau}_L \tau_R}$ according to the vector operator.

Finally, the $R_{K^{(*)}}$ anomaly can be accommodated by setting the remaining independent parameters of the model $f_{\bar{\nu}_L \nu_R} = g_{\bar{\nu}_L \nu_R}$, resulting in $O_9 = -O_{10} = -0.67$, as required by global fits of the anomalous signal.

B. Compatibility of the two anomalies

A possible problem for the simultaneous explanations of the charged and neutral current flavor anomalies arises from the definition of $R_{D^{(*)}}$,

$$R_{D^{(*)}} = \frac{\mathcal{B}(\bar{B} \rightarrow D^{(*)} \tau \bar{\nu})}{\mathcal{B}(\bar{B} \rightarrow D^{(*)} \ell \bar{\nu})}, \quad (37)$$

where the denominator contains an average over muons and electrons. It could be consequently thought that the couplings introduced in Eq. (12) to explain the $R_{K^{(*)}}$ anomaly (in orange), if sizeable, may dilute the yield of the model to $R_{D^{(*)}}$. However, because the contribution of the diagram in the left panel of Fig. 10 is negligible, the analyses of $R_{D^{(*)}}$ and $R_{K^{(*)}}$ are essentially uncorrelated

FIG. 8. Values of the lepton and RHN Yukawa couplings resulting in the benchmark points presented in Fig. 7 for the indicated choice of the quark Yukawa couplings. All points respect the bounds discussed in Sec. II C.

within the present framework. This is due to the fact that the couplings employed in the muon sector to explain the $R_{K^{(*)}}$ anomaly are much smaller than the ones entering

the expression for $R_{D^{(*)}}$, as shown in the right panel of Fig. 10.

C. Impact of new ϵ'/ϵ determinations

To conclude our analysis, we investigate the robustness of our model with respect to another flavor observable that recently received increasing attention [72,73]: ϵ'/ϵ . This quantity, which quantifies the direct CP violation in $K \rightarrow \pi\pi$ decays, has been debated in the literature as a further indication of new physics in flavor measurements. The discussion has been revived after the lattice QCD result obtained by the RBC-UKQCD group [74,75]

$$\text{Re}(\epsilon'/\epsilon) = 1.38(5.15)(4.59) \times 10^{-4} \quad (38)$$

was corroborated by further dedicated analyses [76,77], resulting in independent estimates of the SM contribution that all highlight a 2σ deviation from the current experimental measure [78–80]

$$\text{Re}(\epsilon'/\epsilon)_{\text{exp}} = (16.6 \pm 2.3) \times 10^{-4}. \quad (39)$$

FIG. 9. Values of the quark Yukawa couplings that result in the scalar and pseudoscalar Wilson coefficients indicated in Fig. 7. The plot is obtained by letting the scalar masses vary in the 300–350 GeV range and setting $f_{\bar{\nu}_L \tau_R}$ as indicated by the vector Wilson coefficients.

FIG. 10. Left panel: The diagram that hinders simultaneous explanations for $R_{D^{(*)}}$ and $R_{K^{(*)}}$ in the present framework. Right panel: Values of C_{VL}^μ resulting from the Yukawa couplings used to explain the $R_{K^{(*)}}$ anomaly.

Presently, the origin of the anomaly remains unclear as misdeterminations of the SM contribution alone could explain the mentioned deviation [73]. In light of this ambiguous signal, we discuss below how new determinations of the SM contribution into ϵ'/ϵ could constrain, or be rectified, in the present framework.

The contribution of a charged scalar to the effective $d \rightarrow sg$ Hamiltonian

$$\mathcal{H}_{d \rightarrow sg} = -\frac{G_F}{\sqrt{2}} V_{ts}^* V_{td} (C_{8g} O_{8g} + C'_{8g} O'_{8g}), \quad (40)$$

where the chromomagnetic dipole operators (CMOs) are

$$O_{8g} = \frac{g_s}{8\pi^2} \bar{s} \sigma^{\mu\nu} T^a P_L d G_{\mu\nu}^a, \quad O'_{8g} = \frac{g_s}{8\pi^2} \bar{s} \sigma^{\mu\nu} T^a P_R d G_{\mu\nu}^a, \quad (41)$$

results in a short-distance modification of the SM dynamics which, without introducing further sources of CP violation, alleviates the tension with the current experimental results [81,82].

The determination of the hadronic matrix element for $K^0 \rightarrow \pi\pi$ in dual QCD [83–87] allows an estimate of the

order of magnitude needed for the Wilson coefficients in Eq. (40) to reproduce the ϵ'/ϵ measurement [77]

$$\text{Re}(\epsilon'/\epsilon)_{8g} \sim -(1.85 \times 10^5 \text{ GeV}) \times \text{Im}(C_{8g}^-), \quad (42)$$

where C_{8g}^- is the combination

$$C_{8g}^- = -\frac{G_F}{\sqrt{2}} V_{ts}^* V_{td} (m_d C'_{8g} - m_s C_{8g}). \quad (43)$$

We stress that Eq. (42) is obtained under the assumption of a SM-like value for the indirect CP violation in ϵ_K enforced in the framework at hand by the same bounds targeting the mass splitting and CP violation of the neutral K system [88]. As previously shown in Ref. [31], these constraints do not significantly limit the parameter space selected by the anomalous signals.

With the Yukawa texture in Eq. (13), the same interactions behind our explanations of the $R_{D^{(*)}}$ and $R_{K^{(*)}}$ anomalies yield new contributions to the CMOs that source ϵ'/ϵ . In particular, the propagation of H_2^- in Fig. 11 results in

FIG. 11. Left panel: The dominant 3HDM contribution to ϵ'/ϵ . Right panel: Values of $f_{d_L t_R}$ required to match the measured value $\text{Re}(\epsilon'/\epsilon)_{8g} = 10^{-3}$ for the indicated choices of the quark Yukawa couplings.

$$\begin{aligned}
C_{8g_{H_2}}^- = & \frac{1}{4} V_{tb} V_{ts}^* |g_{\bar{b}_L t_R}|^2 \{m_d (C_{cc12}[m_{H_2}^2, m_t^2, m_t^2] + C_{cc2}[m_{H_2}^2, m_t^2, m_t^2] + C_{cc22}[m_{H_2}^2, m_t^2, m_t^2]) + \\
& - m_s (C_{cc1}[m_{H_2}^2, m_t^2, m_t^2] + C_{cc11}[m_{H_2}^2, m_t^2, m_t^2] + C_{cc12}[m_{H_2}^2, m_t^2, m_t^2])\} \\
& \times V_{tb} V_{ts}^* \frac{m_d - m_s}{48(m_{H_2}^2 - m_t^2)^4} |g_{\bar{b}_L t_R}|^2 (2m_{H_2}^6 + 3m_{H_2}^4 m_t^2 - 6m_{H_2}^2 m_t^4 - 6m_{H_2}^4 m_t^2 \log(m_{H_2}^2/m_t^2) + m_t^6), \quad (44)
\end{aligned}$$

and, in turn, in value of ϵ'/ϵ comparable in magnitude to the current SM estimate. Therefore, should a revised computation of the SM contribution explain the observed discrepancy, precision measurements of this parameter will certainly offer a new way to probe our scenario. On the contrary, should the ϵ'/ϵ measurement require the presence of new physics, our framework could offer a possible solution through a minimal modification of the proposed Yukawa texture

$$\mathcal{Y}_1^d = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & f_{\bar{c}_L b_R} \\ 0 & 0 & 0 \end{pmatrix} \rightarrow \mathcal{Y}_1^d = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & f_{\bar{c}_L b_R} \\ 0 & 0 & f_{\bar{s}_L t_R} \end{pmatrix}, \quad (45)$$

where the new diagonal down-Yukawa coupling $f_{\bar{s}_L t_R}$ results, via the H_1^- contribution shown in Fig. 11, in the term

$$\begin{aligned}
C_{8g_{H_1}}^- = & -\frac{m_t}{4} V_{tb} V_{cs}^* (f_{\bar{c}_L b_R} f_{\bar{s}_L t_R}) (C_{cc1}[m_{H_1}^2, m_t^2, m_t^2] \\
& + C_{cc2}[m_{H_1}^2, m_t^2, m_t^2]) \\
= & V_{tb} V_{cs}^* \frac{m_t}{8(m_{H_1}^2 - m_t^2)^3} (f_{\bar{c}_L b_R} f_{\bar{s}_L t_R}) \\
& \times (3m_{H_1}^4 - 4m_{H_1}^2 m_t^2 - 2m_{H_1}^4 \log(m_{H_1}^2/m_t^2) + m_t^4). \quad (46)
\end{aligned}$$

Notice the relevance of the Yukawa term $f_{\bar{c}_L t_R}$, which is therefore pivotal for the explanation of all of the anomaly discussed in the present work and offers a clear signature for the investigation of our proposal.

In Fig. 11, we assess the value of $f_{\bar{d}_L t_R}$ required to have $Re(\epsilon'/\epsilon)_{8g} = 10^{-3}$, considering the values of $f_{\bar{c}_L t_R}$ and $g_{\bar{b}_L t_R}$ used in the benchmark points of the previous section. Clearly, it is sufficient to set $f_{\bar{d}_L t_R} \sim 10^{-1}$ to enhance the 3HDM contribution into the CMOs and explain the ϵ'/ϵ anomaly on top of the $R_{D^{(*)}}$ and $R_{K^{(*)}}$ signals.

V. DISCUSSION AND CONCLUSIONS

We attempted to demystify the implications of the $R_{D^{(*)}}$ and $R_{K^{(*)}}$ anomalies in the context of physics beyond the standard model. Although these signals seem to indicate new lepton flavor violating physics at very different scales, we demonstrated that this fact does not necessarily imply the existence of exotic new physics. In fact, the full

dynamics responsible for all of the analyzed measurements can still be addressed within a more conventional framework using model building elements already available in the standard model.

Specifically, we showed that the present $R_{D^{(*)}}$ and $R_{K^{(*)}}$ anomalies can be simultaneously explained in a three Higgs doublets model extended with right-handed neutrinos. The results plotted in Fig. 7 demonstrate that the model predicts values of $R_{D^{(*)}}$ within the 1σ region indicated by the experiments, in accordance with the remaining phenomenological constraints and for perturbative values of the involved Yukawa couplings. The result is achieved owing to the interplay between the loop-induced vector operators and tree-level scalar operators in Eq. (4), which depend on different sets of parameters. We furthermore showed that in the considered scheme, the $R_{D^{(*)}}$ and $R_{K^{(*)}}$ anomalies arise from independent interactions of the additional scalar doublets and consequently found a simultaneous explanation for the two measurements.

The robustness of our results was tested against a further flavor observable, ϵ'/ϵ . Once again, we found that within the proposed 3HDM, this anomaly can be modeled independently from the remaining flavor physics signals. Should the current deviation in this parameter be confirmed as an effect of new physics, our proposal would then provide the first framework able to simultaneously explain all of the mentioned anomalies.

Our analyses omitted the angular observables encoded in the P'_i parameters. In this regard, we only remark that P'_5 certainly receives additional contributions in the present model via the C_9 Wilson coefficient. These contributions, however, do not allow us to explain the small tension observed in the measurements of this parameter, which select negative values of C_9 with possible subdominant negative contributions from C_{10} [89]. Because our solution selects $C_9 = -C_{10}$ instead, further studies of angular observables have certainly the potential to test the framework.

A crucial aspect of our work is the inclusion of GeV-scale right-handed neutrinos, which allowed for the sizable loop-level contributions needed to explain the analyzed signals. It is remarkable that the presence of these particles in nature is advocated also in connection to other open problems of contemporary physics [90–92], and our framework therefore connects the $R_{D^{(*)}}$ and $R_{K^{(*)}}$ flavor anomalies to the well-known phenomenology of neutrino masses and to the puzzle of the baryon asymmetry of the

Universe. The considered right-handed neutrinos are also a primary target of the next-generation beam-dump experiment SHIP at CERN, which therefore has the potential to directly test our scenario on top of complementary collider searches for new scalar particles.

In light of the proposed scheme, it is natural that the first signals of new physics would be detected in low-energy flavor experiments rather than in high-energy collider searches and precision observables: Because the GeV-scale right-handed neutrinos phenomenology is intrinsically a low-energy phenomenon, the first signals of their presence naturally occur in the loop processes behind the flavor observables under discussion, preceding the potential collider signatures of still undiscovered scalar particles.

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Note added.—Recently, the LHCb Collaboration presented new results concerning the R_K anomaly at the 2019 Moriond EW Conference [93]. The new measurement is based on both a reanalysis of the LHC run I data finding

$$R_K^{\text{rev run I}} = 0.717^{+0.083}_{-0.071} \pm 0.017, \quad (47)$$

as well as on part of the run II data

$$R_K^{\text{run II}} = 0.928^{+0.089}_{-0.076} \pm 0.020, \quad (48)$$

and is compatible with both the SM expectation and the previous result. Accounting for the new data yields an average $R_K = 0.846^{+0.060}_{-0.054} \pm 0.015$ which falls 2.5 standard deviations below unity. In addition, the Belle Collaboration reported their preliminary results on R_{K^*} obtained by averaging over observations of B^0 and B^+ decays. These indicate [94]

$$R_{K^*} = \begin{cases} 0.52^{+0.36}_{-0.26} \pm 0.05 & \text{for } 0.045 \text{ GeV}^2 \leq q^2 \leq 1.1 \text{ GeV}^2, \\ 0.96^{+0.45}_{-0.29} \pm 0.11 & \text{for } 1.1 \text{ GeV}^2 \leq q^2 \leq 6 \text{ GeV}^2, \end{cases} \quad (49)$$

where q^2 is the invariant mass of the final state dilepton system.

Whereas the new result shows a global trend that is closer to the SM predictions, improvement in the sensitivity of the observations has reduced the experimental errors in a way that the total significance of the anomalies is only marginally impaired. From the point of view of our solution, explaining these new results would then simply require a slight reduction of the Yukawa couplings entering the $b \rightarrow s$ current.

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