Right-handed neutrinos and $R(D^{(*)})$

Dean Robinson,\textsuperscript{a} Bibhushan Shakya\textsuperscript{a,b} and Jure Zupan\textsuperscript{a}

\textsuperscript{a}Department of Physics, University of Cincinnati, Cincinnati, Ohio 45221, U.S.A.
\textsuperscript{b}Leinweber Center for Theoretical Physics (LCTP), University of Michigan, Ann Arbor, Michigan 48109, U.S.A.

E-mail: dean.robinson@uc.edu, shakyabn@ucmail.uc.edu, zupanje@ucmail.uc.edu

ABSTRACT: We explore scenarios where the $R(D^{(*)})$ anomalies arise from semitauonic decays to a right-handed sterile neutrino. We perform an EFT study of all five simplified models capable of generating at tree-level the lowest dimension electroweak operators that give rise to this decay. We analyze their compatibility with current $R(D^{(*)})$ data and other relevant hadronic branching ratios, and show that one simplified model is excluded by this analysis. The remainder are compatible with collider constraints on the mediator semileptonic branching ratios, provided the mediator mass is of order TeV. We also discuss the phenomenology of the sterile neutrino itself, which includes possibilities for displaced decays at colliders and direct searches, measurable dark radiation, and gamma ray signals.

KEYWORDS: Beyond Standard Model, Heavy Quark Physics, Neutrino Physics

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1 Introduction

Measurements of the semitauonic to light semileptonic ratios at multiple experiments [1–6],

\[ R(D^{(*)}) = \frac{\text{Br}[B \to D^{(*)}\tau\bar{\nu}]}{\text{Br}[B \to D^{(*)}\ell\bar{\nu}]}, \quad l = e, \mu, \]

(1.1)

exhibit a 4σ tension with respect to the Standard Model (SM) predictions, once both \( D \) and \( D^* \) measurements are combined [7] (see also refs. [8–12]). Beyond the Standard Model (BSM) explanations of this anomaly typically require new physics (NP) close to the TeV scale. Since the SM neutrino is part of an electroweak doublet, corresponding constraints necessarily arise from high-\( p_T \) measurements of \( pp \to \tau^+\tau^- \) at the LHC [13], \( Z \) and \( \tau \) decays [14, 15], and contributions to flavor changing neutral currents (FCNCs), that can be severe.

As discussed in refs. [16, 17] (see also refs. [18, 19]), the observed enhancements of \( R(D^{(*)}) \) can be achieved not only through NP contributions to the \( b \to c\tau\bar{\nu}_\tau \) decay, where
is the SM left-handed \( \tau \) neutrino, but also via a new decay channel, \( b \to c \tau N_R \), where \( N_R \) is a sterile right-handed neutrino. The \( b \to c \tau \bar{\nu} \) decay becomes an incoherent sum of two contributions: to streamline notation we denote \( \nu = N_R \) or \( \nu_\tau \), so that \( \text{Br}[b \to c \tau \bar{\nu}] = \text{Br}[b \to c \tau \bar{\nu}_\tau] + \text{Br}[b \to c \tau N_R] \). Since the NP couples to right-handed neutrinos, this can relax many of the electroweak constraints from the \( \tau \) processes mentioned above.

In the specific context of refs. [16, 17], the \( b \to c \tau N_R \) decay is mediated by an SU(2)\(_L\) singlet \( W^0 \), which can be UV completed in a ‘3221’ model. In this paper we generalize the EFT studies of refs. [16, 17] to the full set of dimension-six operators involving \( N_R \) (for earlier partial studies see [20–22]). Assuming that the NP corrections are due to a tree level exchange of a new mediator, there are five possible simplified models for \( b \to c \tau N_R \), whose mediators are: the SU(2)\(_L\)-singlet vector boson — the \( W^0 \); a scalar electroweak doublet; and three leptoquarks.

For each simplified model we identify which regions of parameter space are consistent with the \( R(D^{(*)}) \) anomaly, subject to exclusions from the \( B_c \to \tau \nu \) branching ratio [23–25]. We further examine the variation in the signal differential distributions expected for each simplified model. While some electroweak constraints are relaxed, these simplified models nonetheless typically imply various sizeable semileptonic branching ratios for the tree-level mediators, for which moderately stringent collider bounds already exist. We show that, depending on the ratios of NP couplings in the simplified model, these in turn set lower bounds of \( \mathcal{O}(\text{TeV}) \) on the mediator masses. We then proceed to examine the implications for neutrino phenomenology, such as bounds from radiative contributions to the SM neutrino masses, astrophysical constraints from sterile neutrino electromagnetic decays, plausible cosmological histories that admit these sterile neutrinos, and displaced decays at colliders and direct searches. In our analysis, we will require the \( N_R \) to be light — \( m_{N_R} \lesssim \mathcal{O}(100) \) MeV — in order not to disrupt the measured missing invariant mass spectrum in the full \( \mathcal{B} \to D^{(*)} \tau \bar{\nu} \) decay chain. Whether heavier sterile neutrinos are compatible with data requires a dedicated forward-folded study, performed by the experimental collaborations.

The paper is structured as follows. Section 2 contains the EFT analysis of the \( R(D^{(*)}) \) data for the case of the right-handed neutrino and introduces the five possible tree-level mediators. Collider constraints on these simplified models are studied in section 3, while section 4 contains the related sterile neutrino phenomenology. Our conclusions follow in section 5. Appendix A examines the structure of the \( b \to c \tau \bar{\nu} \) differential distributions for the simplified models.

# 2 EFT analysis

## 2.1 EFTs and simplified models

We consider the extension of the SM field content by a single new state, a right handed, sterile neutrino transforming as \( N_R \sim (1, 1, 0) \) under SU(3)\(_c\) × SU(2)\(_L\) × U(1)\(_Y\). This state may couple to the SM quarks via higher dimensional operators. Above the electroweak scale, one therefore adds to the renormalizable SM Lagrangian the following effective in-
tions,
\[ \mathcal{L}^{\text{EW}}_{\text{eff}} = \sum_{a,d} \frac{C_{ad}}{\Lambda_{\text{eff}}^{d-4}} Q_a + \cdots, \quad (2.1) \]
where \(Q_a\) are dimension-\(d\) operators, \(C_{ad}\) are the corresponding dimensionless Wilson coefficients (WCs), and \(\Lambda_{\text{eff}}\) is the effective scale defined to be
\[ \Lambda_{\text{eff}} = (2\sqrt{2} G_F V_{cb})^{-1/2} \simeq 0.87 \left(\frac{40 \times 10^{-3}}{V_{cb}}\right)^{1/2} \text{TeV}. \quad (2.2) \]

The most general basis of dimension-6 operators that can generate the charged current \(b \to c\tau\bar{N}_R\) decay is given by
\[
Q_{SR} = \epsilon_{ab} (\bar{Q}_L^a d_R) (\bar{L}_R^b N_R), \quad Q_{SL} = (\bar{u}_R Q_L^a) (\bar{L}_R^b N_R), \quad (2.3a)
\]
\[
Q_T = \epsilon_{ab} (\bar{Q}_L^a \sigma^{\mu \nu} d_R) (\bar{L}_L^b \sigma_{\mu \nu} N_R), \quad Q_{VR} = (\bar{u}_R \gamma^\mu d_R) (\bar{e}_R \gamma_\mu N_R). \quad (2.3b)
\]
Here \(a, b\) are SU(2)_L indices, \(\epsilon_{ab}\) is an antisymmetric tensor with \(\epsilon_{12} = -\epsilon_{21} = 1\), and we use the four-component notation, with \(Q_L\) the SM quark doublet, \(u_R\) and \(d_R\) the up- and down-quark singlets, and \(L_L\) the SM lepton doublet. (As usual, there is only one non-vanishing tensor operator, since \(\sigma_{\mu \nu} \otimes \sigma^{\mu \nu} = 0\), which immediately follows from the relation \(\sigma_{\mu \nu} \otimes \sigma^{\mu \nu} \gamma_5 = \sigma_{\mu \nu} \gamma_5 \otimes \sigma^{\mu \nu}\).) One may also include the dimension-8 operator
\[ Q_{VL} = (\bar{Q}_L \tilde{H} \gamma^\mu H^\dagger Q_L) (\bar{\ell}_R \gamma_\mu N_R), \quad (2.4) \]
where \(\tilde{H} = \epsilon H^\ast\), as well as the operators with the left-handed sterile neutrino field, \(N_{R}\), that start at dimension-7,
\[
Q_{SR}' = (\bar{Q}_L \tilde{H} d_R) (\bar{\ell}_R N_R^c), \quad Q_{SL}' = (\bar{u}_R H^\dagger Q_L) (\bar{\ell}_R N_R^c), \quad (2.5a)
\]
\[
Q_T' = (\bar{u}_R \sigma^{\mu \nu} H^\dagger Q_L) (\bar{\ell}_R \sigma_{\mu \nu} N_R^c), \quad Q_{VR}' = (\bar{u}_R \gamma^\mu d_R) (\bar{L}_L H \gamma_\mu N_R^c), \quad (2.5b)
\]
and the dimension-9 equivalent of \(Q_{VL}\),
\[ Q_{VL}' = (\bar{Q}_L \tilde{H} \gamma^\mu H^\dagger Q_L) (\bar{L}_L H \gamma_\mu N_R^c). \quad (2.6) \]
Each of the SM fields also carries a family index, i.e., \(Q_{L}, u_R, d_R, L_L\), \(i = 1, 2, 3\), and similarly for the Wilson coefficients, \(C^{ijk}_{ad}\), and the operators, \(Q_{ijk}\), in eq. (2.1), which we have omitted for the sake of simplicity. Since we focus exclusively on the generation of \(b \to c\tau\bar{\nu}\) decays below, we drop the family indices hereafter, unless otherwise stated. Consistency with bounds from direct searches requires that the Wilson coefficients in eq. (2.1) be at most \(O(1)\).

Below the electroweak scale, the top quark, the Higgs, and the \(W\) and \(Z\) bosons are integrated out. At the scale \(\mu \sim m_{c,b}\), the effective Lagrangian, including SM terms (see, e.g., [26]), can be written
\[ \mathcal{L}_{\text{eff}} = \mathcal{L}^{\text{SM}}_{\text{eff}} + \frac{1}{\Lambda_{\text{eff}}^2} \sum_i c_i \mathcal{O}_i, \quad (2.7) \]
in which the NP contributions to $b \to c\tau\bar{\nu}$, induced by the dimension-6 operators in (2.3), are described by the following four-fermion operators,

\[ \mathcal{O}_{SR} = (\bar{c}_L b_R)(\bar{\tau}_LN_R), \quad \mathcal{O}_{SL} = (\bar{c}_R b_L)(\bar{\tau}_LN_R), \]
\[ \mathcal{O}_{VR} = (\bar{c}_R \gamma^\mu b_R)(\bar{\tau}_R \gamma^\mu N_R), \quad \mathcal{O}_{T} = (\bar{c}_L \sigma^{\mu\nu} b_R)(\bar{\tau}_L \sigma_{\mu\nu} N_R). \] (2.8a)

The scalar and tensor operators run under the Renormalization Group. The RG evolution from $M > m_t$ to $\mu < m_b$ gives at one-loop order in the leading log approximation for the Wilson coefficients at the low scale [27, 28], for $X = \text{SR, SL, T},$

\[
c_X(\mu) = \left( \frac{\alpha(m_b)}{\alpha(\mu)} \right)^{\gamma_X/2\beta_0^{(n)}} \left[ \frac{\alpha(m_t)}{\alpha(m_b)} \right]^{\gamma_X/2\beta_0^{(5)}} \frac{\alpha(M)}{\alpha(m_t)} \frac{\gamma_X/2\beta_0^{(6)}}{c_X(M)} \equiv \rho_X(\mu; M)c_X(M), \] (2.9)

with anomalous dimensions $\gamma_{\text{SR,SL}} = -8$, $\gamma_T = 8/3$ and the one loop $\beta$-function coefficient $\beta_0^{(n)} = 11 - 2n/3$. The running of $c_{\text{SR,SL,T}}$ depends only weakly on the high scale $M$, and hereafter we set $M = \Lambda_{\text{eff}}$. Fixing the scale low scale to $\mu = \sqrt{m_b m_c}$ — anticipating the chosen matching scale of QCD onto HQET for the $B \to D^{(*)}$ form factor parametrization — one finds

\[
\rho_{\text{SR,SL}} \simeq 1.7, \quad \rho_T \simeq 0.84 . \] (2.10)

Assuming the flavor indices are given in the mass eigenstate basis, the NP operators (2.1) can be matched onto the operators (2.3) as $c_X(\Lambda_{\text{eff}}) = C_X^{233}$, neglecting the tiny mixing of active neutrinos into $N_R$. Note that the operators $\mathcal{O}_{\text{SR,T,SL}}$ are accompanied by the SU(2)$_L$ related operators

\[
\mathcal{O}_{\text{SR}} = (\bar{s}_L b_R)(\bar{\nu}_e N_R), \quad \mathcal{O}_{T}^e = (\bar{s}_L \sigma^{\mu\nu} b_R)(\bar{\nu}_e \sigma_{\mu\nu} N_R), \] (2.11)

and $(\bar{c}_R t_L)(\bar{\nu}_e N_R)$. The Wilson coefficients of these operators, $c_{\text{SR,T,SL}}^e$, correspond to $c_{\text{SR,T,SL}}$, respectively, up to one-loop or higher-order corrections.

Each of the dimension-six operators in eq. (2.3) can arise from the tree level exchange of a new state, either a scalar or a vector. The possible mediators, together with the Wilson coefficients $c_X$ they can contribute to, are listed in table 1. Two of these mediators are color singlets: the charged vector resonance $W'_\mu$, discussed extensively in refs. [16, 17], and the weak doublet scalar $\Phi$. The remaining mediators are leptoquarks, for which we use the notation from ref. [29]. In some cases the structure of the mediator Lagrangian, $\mathcal{L}_{\text{int}}$, implies relations between the various Wilson coefficients, denoted by equalities in table 1. In particular, for the $R_2$ and $S_1$ models, $c_{\text{SR}}(\Lambda_{\text{eff}}) = \pm 4c_T(\Lambda_{\text{eff}})$, which evolves to

\[
c_{\text{SR}}(\mu) = \pm 4c_T(\mu), \quad r = \rho_{\text{SR}}/\rho_T \simeq 2.0 , \] (2.12)

at the $B$ meson scale.

For completeness, we list the remaining $b \to c\tau\bar{N}_R$ dimension-6 operators at $\mu \sim m_{c,b}$,

\[
\mathcal{O}_{\text{SR}}' = (\bar{c}_L b_R)(\bar{\tau}_R N_R^c), \quad \mathcal{O}_{\text{SL}}' = (\bar{c}_R b_L)(\bar{\tau}_R N_R^c), \]
\[ \mathcal{O}_{\text{VR}}' = (\bar{c}_R \gamma^\mu b_R)(\bar{\tau}_R \gamma^\mu N_R^c), \quad \mathcal{O}_{\text{VL}}' = (\bar{c}_L \gamma^\mu b_R)(\bar{\tau}_L \gamma^\mu N_R^c), \]
\[ \mathcal{O}_{T}' = (\bar{c}_R \gamma^{\mu\nu} b_R)(\bar{\tau}_R \gamma^{\mu\nu} N_R^c), \quad \mathcal{O}_{\text{VL}}' = (\bar{c}_L \gamma^{\mu\nu} b_R)(\bar{\tau}_L \gamma^{\mu\nu} N_R^c). \] (2.13a)
in the heavy quark expansion, with matching scale $L$ the form factor $t_\ell$ for NP corrections to $R$ regions in the $c$ an incoherent sum of two contributions: the SM decay $b$ With the addition of a right-handed neutrino decay mode, the $N_R$ in ref. [8] (see also refs. [9, 10]), are The SM predictions, e.g. making use of the model-independent form factor $t_\ell$ $L$ The present experimental world-averages for $R(D^{(*)})$ data The SM predictions, e.g. making use of the model-independent form factor fit ' $L_w$' of ref. [8] (see also refs. [9, 10]), are With the addition of a right-handed neutrino decay mode, the $B \rightarrow D^{(*)}\tau\bar{\nu}$ decays become an incoherent sum of two contributions: the SM decay $b \rightarrow c\tau\bar{\nu}_\tau$ and the new mode $b \rightarrow c\tau N_R$. The $N_R$ contributions therefore increase both of the $B \rightarrow D^{(*)}\tau\bar{\nu}$ branching ratios above the SM predictions, as would be required to explain the experimental measurements of $R(D^{(*)})$. In figure 1, we show for each simplified model of table 1 the accessible contours or regions in the $R(D) - R(D^*)$ plane, compared to the experimental data. The predictions for NP corrections to $R(D^{(*)})$ are obtained from the expressions in ref. [30], making use of the form factor fit ‘ $L_w$’ of ref. [8]. This fit was performed at next-to-leading order in the heavy quark expansion, with matching scale $\mu = \sqrt{m_\phi m_c}$ and quark masses defined

| mediator | irrep | $\delta L_{int}$ | WCs |
|----------|------|-----------------|-----|
| $W'_\mu$ | $(1,1)_1$ | $g' (c_q \bar{u}_R \gamma_\mu d_R + \frac{c_N}{\Lambda} \bar{\ell}_R \gamma_\mu N_R) W'_\mu$ | $c_{VR}$ |
| $\Phi$ | $(1,2)_{1/2}$ | $y_u \bar{u}_R Q_L \ell \Phi + y_d \bar{d}_R Q_L \ell \Phi + y_N \bar{N}_R L_L \ell \Phi$ | $c_{SL}(\mu), c_{SR}(\mu)$ |
| $U^\mu_1$ | $(3,1)_{2/3}$ | $(\alpha_{LQ} \bar{L}_L \gamma_\mu Q_L + \alpha_{d\ell} \bar{\ell}_R \gamma_\mu d_R) U^\mu_1 + \alpha_{uN} (\bar{u}_R \gamma_\mu N_R) U^\mu_1$ | $c_{SL}(\mu), c_{VR}$ |
| $\tilde{R}_2$ | $(3,2)_{1/6}$ | $\alpha_{Ld} (\bar{L}_L d_R) \ell \tilde{R}_2 + \alpha_{QN} (\bar{Q}_L N_R) \tilde{R}_2$ | $c_{SR}(\mu) = 4 r_{CT}(\mu)$ |
| $S_1$ | $(3,1)_{1/3}$ | $z_u (\bar{U}_R \ell_R) S_1 + z_d (\bar{d}_R N_R) S_1 + z_Q (\bar{Q}_L \ell L_L) S_1$ | $c_{VR}, c_{SR}(\mu) = -4 r_{CT}(\mu)$ |

Table 1. The tree-level mediators that can generate the four-fermion operators with right-handed neutrino, $N_R$, in eqs. (2.8). The relevant Wilson coefficients are shown in the final column, explicitly defined at scale $\mu$ where relevant, and including the factor $r \equiv \rho_{SR}/\rho_T \simeq 2.0$. The generation of these operators from the electroweak scale four-Fermi operators (2.4)–(2.6) requires additional insertions of the Higgs vev, $v_{EW}$, and, apart from $\mathcal{O}_{VL}$, also the left-handed sterile neutrino $N^c_R$. These $\mathcal{O}'_a$ operators are the same as those in ref. [27], but with $N^c_R$ replacing the SM neutrino $\nu$. Eqs. (2.8) and (2.13) together form a complete basis of $b \rightarrow c\tau N_R$ dimension-six four-fermion operators. Since the Wilson coefficients of the operators in eq. (2.13) are suppressed by additional powers of $v_{EW}/\Lambda_{eff}$, we will only focus on the dimension-6 operators listed in eq. (2.3) and (2.8) in the remainder of this paper.

2.2 Fits to $R(D^{(*)})$ data

The present experimental world-averages for $R(D^{(*)})$ are [7]

$$R(D)|_{exp} = 0.407 \pm 0.046, \quad R(D^*)|_{exp} = 0.304 \pm 0.015, \quad \text{corr.} = -0.20.$$ (2.14)

The SM predictions, e.g. making use of the model-independent form factor fit ‘$L_{w \geq 1}$SR’ of ref. [8] (see also refs. [9, 10]), are

$$R(D)|_{th} = 0.299 \pm 0.003, \quad R(D^*)|_{th} = 0.257 \pm 0.003, \quad \text{corr.} = +0.44.$$ (2.15)

With the addition of a right-handed neutrino decay mode, the $B \rightarrow D^{(*)}\tau\bar{\nu}$ decays become an incoherent sum of two contributions: the SM decay $b \rightarrow c\tau\bar{\nu}_\tau$ and the new mode $b \rightarrow c\tau N_R$. The $N_R$ contributions therefore increase both of the $B \rightarrow D^{(*)}\tau\bar{\nu}$ branching ratios above the SM predictions, as would be required to explain the experimental measurements of $R(D^{(*)})$. In figure 1, we show for each simplified model of table 1 the accessible contours or regions in the $R(D) - R(D^*)$ plane, compared to the experimental data. The predictions for NP corrections to $R(D^{(*)})$ are obtained from the expressions in ref. [30], making use of the form factor fit ‘$L_{w \geq 1}$SR’ of ref. [8]. This fit was performed at next-to-leading order in the heavy quark expansion, with matching scale $\mu = \sqrt{m_\phi m_c}$ and quark masses defined
Figure 1. The enhancements of $R(D^{(*)})$ from $b \rightarrow c \tau \bar{N}_R$ decays for various simplified models. The world average experimental 1σ, 2σ, and 3σ fit regions are shown in decreasing shade of gray. The SM point is denoted by a black dot.

in the $\Upsilon(1S)$ scheme, relevant for a self-consistent treatment of the $B_c \rightarrow \tau \nu$ constraints below. The $W'$ and $\tilde{R}_2$ simplified models have only a single free Wilson coefficient and are constrained to a contour: since the $N_R$ contributions add incoherently to the SM, the phase of each Wilson coefficient is unphysical. By contrast, $\Phi$, $U_1$, and $S_1$ have two free Wilson coefficients, corresponding to two free magnitudes and a physical relative phase, permitting them to span a region.

Assuming first that all Wilson coefficients are real, we show in figure 2 the 0.5σ, 1σ CLs (dark, light blue) and 1.5σ, 2σ CLs (dark, light green) in the relevant Wilson coefficient spaces for each simplified model. These CLs are generated by the $\chi^2$ defined with respect to the $R(D^{(*)})$ experimental data and correlations (2.14), not including the possible effects of NP errors. That is,

$$\chi^2 = v^T \sigma^{-1}_{R(D^{(*)})} v, \quad v = (R(D)_{\text{th}} - R(D)_{\text{exp}}, R(D^{*})_{\text{th}} - R(D^{*})_{\text{exp}}),$$

(2.16)

The $\chi^2$ CLs (dof = 2) in figure 2 then correspond simply to projections of the CL ellipses in figure 1. We will hereafter refer to the minimal $\chi^2$ points in the WC space for each simplified model as the model’s ‘best fit’ points with respect to the $R(D^{(*)})$ results (2.14), though it should be emphasized that this is not the same as a NP WC fit to the experimental data, which would require inclusion of the NP errors in the underlying experimental fits. In figure 2 the best fit points are shown by black dots, with explicit values provided in table 2. For the $W'$ and $\tilde{R}_2$ models, we show the explicit $\chi^2$, as well as the intervals corresponding to 1σ and 2σ CLs (dof = 2).

The additional NP currents from the operators (2.8) also incoherently modify the $B_c \rightarrow \tau \nu$ decay rate with respect to the SM contribution (cf. refs. [23, 24]), such that

$$\text{Br}(B_c \rightarrow \tau \nu) = \frac{\tau_{B_c} f_{B_c}^2 m_{B_c}^2}{64 \pi \Lambda^4_{\text{eff}}} \left(1 - m_{\tau}^2 / m_{B_c}^2\right) \left[1 + |c_{\text{VR}} + \frac{m_{B_c}^2 c_{\text{SL}}^{(\mu)} - c_{\text{SR}}^{(\mu)}}{m_{\tau} (m_b + m_c)}|^2\right]^2,$$

(2.17)
ms quark masses, 

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\[ \chi \]

\[ \gamma \]

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have non-trivial solutions
\[
\cos(\varphi_0) = \begin{cases} 
\frac{0.24-0.51|c_{SR}|^2-0.51|c_{SL}|^2}{|c_{SR}|^2|c_{SL}|}, & \Phi, \\
\frac{0.38-1.38|c_{VR}|^2-0.60|c_{SR}|^2}{|c_{VR}|^2|c_{SR}|}, & U_1, \\
\frac{0.32-1.40|c_{VR}|^2-0.61|c_{SR}|^2}{|c_{VR}|^2|c_{SR}|}, & S_1,
\end{cases}
\]
valid only on the domain $|\cos(\varphi_0)| < 1$, and otherwise $\cos(\varphi_0) = \pm 1$. These phase-optimized CLs for the $\Phi$, $U_1$, and $S_1$ models are shown in figure 3, with the explicit best fit points listed in table 2. The best fit points for $U_1$ and $S_1$ remain the same, and one sees that these models continue to have non-excluded 1$\sigma$ CLs. An additional best fit point emerges for the $\Phi$ simplified model; however, this model remains excluded, and we therefore do not consider it further in this paper.

Finally, the exchange of mediators that generates the $c_{SR,T}$ Wilson coefficients also results in $c_{SR,T}^4$ of similar size (see eq. (2.11)). The two operators in eq. (2.11) contribute to $b \to s\nu\bar{\nu}$ rates. This gives, for instance, for the $B \to K\nu\bar{\nu}$ decay rate (far enough from the kinematic threshold so that we can neglect all the final state masses) [33, 34]

\[
\frac{d\Gamma_{B\to K\nu\bar{\nu}}}{dz}/\frac{d\Gamma_{B\to K\nu\bar{\nu}}}{dz}_{\text{SM}} = 1 + z \frac{32\pi^2}{3\alpha^2} \left| \frac{V_{cb}}{V_{ub} V_{ts}} \right|^2 \left[ \frac{3 (c_{SR}^\epsilon)^2}{8 (1-z)^2} f_0^2 + (c_{T}^\epsilon)^2 f_T^2 \right] \\
\simeq 1 + 5 \times 10^4 z \left[ \frac{3 (c_{SR}^\epsilon)^2}{8 (1-z)^2} f_0^2 + (c_{T}^\epsilon)^2 f_T^2 \right] 
\]

### Table 2

| Model | WC | Real Best fit $\chi^2$ | Phase-optimized Best fit $\chi^2$ |
|-------|----|------------------------|----------------------------------|
| $W'$ | $c_{VR}$ | $\pm 0.46$ 1.0 | - - |
| $\bar{R}_2$ | $(\mu) c_{SR} = 4\tau c_{TR}^\mu$ | $\pm 0.72$ 0.5 | - - |
| $\Phi$ | $\{c_{SR}, c_{TL}^\mu\}$ | $\{\pm 1.50, 0.84\}$ 0. | $\{1.50, -0.84\}$ 0. |
| $\Phi$ | $\{c_{SR}, c_{TL}^\mu\}$ | $\{\pm 0.84, 0.15\}$ 0. | $\{0.84, -1.50\}$ 0. |
| $U_1$ | $\{c_{VR}, c_{SR}^\mu\}$ | $\{\pm 0.45, 0.93\}$ 0. | $\{0.45, -0.93\}$ 0. |
| $S_1$ | $c_{SR} = 4\tau c_{TR}^\mu$ | $\{\pm 0.40, 0.85\}$ 0. | $\{0.40, -0.85\}$ 0. |

Table 2. Best fit points for each model with respect to the $R(D^*)$ results (2.14), for real and phase-optimized Wilson coefficients. In the phase-optimized case, we show best fits up to an overall phase, by choosing the first WC to be real and positive definite.
with the three $B \to K$ form factors, $f_0(q^2)$, $f_+(q^2)$, $f_T(q^2)$, functions of $q^2$, the invariant mass squared of the neutrino pair, and $z = q^2/m_B^2$. The present experimental bound, $\text{Br}(B^+ \to K^+\nu\bar{\nu}) < 1.6 \times 10^{-5}$ [35], is only a factor of a few above the SM prediction, $\text{Br}(B^+ \to K^+\nu\bar{\nu})|_{\text{SM}} \approx 4 \times 10^{-6}$ [36]. This implies that $c_{\text{SR}}^2$ and $c_T^2$ are highly suppressed, to the level of $O(10^{-2})$, introducing tensions with the required size of $c_{\text{SR}}, c_T$ to explain the $R(D^{(*)})$ anomaly. In the single mediator exchange models in table 1, this means that the product $\alpha_{Ld}^2\alpha_{QN}^2$ for $\tilde{R}_2$ and the product $z_Q^{23}z_Q^{23}$ for $S_1$ (and $y_d^{32}$ for $\Phi$) need to be much smaller than what is required to explain $R(D^{(*)})$. This excludes the $\tilde{R}_2$ as a simple one mediator solution to $R(D^{(*)})$: additional operators coupling to the second generation of quark doublets must be introduced, whose couplings are tuned appropriately to suppress the contributions to $b \to sv\nu\bar{\nu}$. However, this approach would in turn induce large radiative contributions to the neutrino masses, which would also need to be tuned away (see section 4). The $S_1$ model also generates too large a $b \to sv\nu$ transition rate at the (non-excluded) best fit point, where $c_{\text{SR}}$ and $c_T$ are nonzero. The dangerous $b \to sv\nu$ contribution can be suppressed by taking $z_Q^{23} \to 0$ (see table 1), which forces $c_{\text{SR}} = c_T \to 0$. This $c_{\text{SR}} = c_T = 0$ point leads to only a small change in $\chi^2$, corresponding to a less than 0.5 $\sigma$ shift in significance, see figure 2.

2.3 Differential distributions

The reliability of the above $R(D^{(*)})$ fit results turns upon the underlying assumption that the differential distributions, and hence experimental acceptances, of the $\bar{B} \to D^{(*)}\tau\nu\bar{\nu}$ decays are not significantly modified in the presence of the NP currents. The $\bar{B} \to D^{(*)}\tau\nu\bar{\nu}$ branching ratios are extracted from a simultaneous float of background and signal data, so that significant modification of the acceptances versus the SM template may alter the extracted values.

To estimate the size of these potential effects, we examine the cascades $\bar{B} \to (D^* \to D\pi)(\tau \to \ell\nu\ell\nu)\bar{\nu}$ and $\bar{B} \to D(\tau \to \ell\nu\ell\nu)\nu$, comparing the purely SM predictions with the predictions for the $2\sigma$ fit regions of the simplified models. We take $N_R$ to be massless, and
include the phase space cuts,
\[ q^2 = (p_B - p_D(\ell))^2 > 4 \text{ GeV}^2, \quad E_\ell > 400 \text{ MeV}, \quad m_{\text{miss}}^2 > 1.5 \text{ GeV}^2, \quad (2.20) \]
as an approximate simulation of the BaBar and Belle measurements performed in refs. [2, 3].
These distributions are generated as in ref. [30], using a preliminary version of the Hammer
library [37]. In appendix A we show the variation of the normalized differential distributions
over the 2t regions in figure 2 — i.e. assuming real couplings, for simplicity — for the
detector observables \( E_D, E_\ell, m_{\text{miss}}^2, \cos \theta_{D\ell} \) and \( q^2 \) compared to the SM distributions.

As already found in ref. [17], the variation of the \( W' \) model with respect to the SM
is negligible. However, the \( \bar{R}_2, U_1 \) and \( S_1 \) theories, since they include interfering scalar
and/or tensor currents, may significantly modify the spectra, as seen also in ref. [30] for
the NP tensor current coupling to a SM neutrino. Thus, a fully self-consistent \( R(D^{(*)}) \) fit
for these models will require a forward-folded analysis by the experimental collaborations:
our analysis above and CLs should be taken only as an approximate guide, within likely
1σ variations in the values of \( R(D^{(*)}) \).

3 Collider constraints on simplified models

The simplified models are subject to low energy flavor constraints as well as bounds from
collider searches. These depend crucially on the assumed flavor structure of the couplings
in table 1. Furthermore, the sensitivity of the collider searches depend on other open decay
channels of the mediators. In this section, we discuss these constraints for the simplified
models.

For the \( S_1 \) and \( \bar{R}_2 \) models, the best fit points are naively excluded by bounds on
\( b \to s \nu \bar{\nu} \) transitions. These can be avoided by including higher dimensional operators, due
to a new set of heavy states, inevitably introducing greater model dependence for LHC
studies. To remain as model independent as possible, we study the collider signatures for
these models using their (\( B_c \to \tau \nu \) consistent) best fit points for \( R(D^{(*)}) \) as a benchmark,
assuming that any new fields required to ameliorate large \( b \to s \nu \bar{\nu} \) (and/or large neutrino
mass contributions) are sufficiently heavy that they do not affect mediator production or
decay.

3.1 \( W' \) coupling to right-handed SM fermions

The charged vector boson \( W'_\mu \) couples to \( SU(2)_L \) singlets only, and transforms as \( W'_\mu \sim (1,1)_1 \), with
\[ \mathcal{L} = \frac{g_V}{\sqrt{2}} c^{ij}_q \bar{u}_R W'_\mu d^j_R + \frac{g_V}{\sqrt{2}} c^i_N \bar{R}_R W'_\mu N_R + \text{h.c.}, \quad (3.1) \]
where \( i, j = 1, 2, 3 \) are generational indices. As in table 1, the coefficients \( c^{ij}_q \) and \( c^i_N \) encode
the flavor structure of the interactions, while \( g_V \) is the overall coupling strength (in simple
gauge models for \( W' \) it can be identified with the gauge coupling constant [16, 17]). A tree
level exchange of \( W' \) generates the operator \( \mathcal{O}_{\text{VR}} \), cf. eqs. (2.8b) and (2.7), with
\[ \frac{c_{\text{VR}}}{\Lambda_{\text{eff}}^2} = -\frac{g^2_{W'} c^{23}_q c^3_N}{2m_{W'}^2}, \quad (3.2) \]
Figure 4. The bound on $\text{Br}(W' \to \tau \nu)$ as a function of $W'$ mass from the 13 TeV ATLAS [38] (solid blue) and CMS [39] (solid red) searches, as well as the projected reach at the end of the high-luminosity LHC run (dashed blue), for the case $c_{q}^{23} = c_{N}^{3}$, $W'$ mass given by eq. (3.3) to fit to $R(D^{(*)})$ data, and the $W'$ couplings to all the other SM quarks set to zero. In this case $\text{Br}(W' \to \tau \nu) = 0.25$ (dashed grey line) if no other $W'$ decay channels are open. All the bounds assume narrow width for $W'$. The region excluded by unitarity is shaded in grey.

The best fit values for $c_{\nu R}$ in table 2 then imply [17]

$$m_{W'} \simeq 540 \left| c_{q}^{23} c_{N}^{3} \right|^{1/2} \left[ \frac{g_{V}}{0.6} \right] \left[ \frac{40 \times 10^{-3}}{V_{cb}} \right]^{1/2} \text{GeV}.$$  (3.3)

In figure 4 we show the minimal set of experimental constraints on such models, applicable to the simplified $W'$ model. For this plot we set $c_{q}^{23} = c_{N}^{3}$, take eq. (3.3) to provide the $W'$ mass that fits the $R(D^{(*)})$ data, and set the $W'$ couplings to all other SM quarks to zero. For this scenario, the ATLAS search at 13 TeV with 36.1 fb$^{-1}$ luminosity [38] and the CMS search with 35.9 fb$^{-1}$ [39] convert to a 95% CL bounds on $\text{Br}(W' \to \tau \nu)$ shown in figure 4 (blue and red lines, respectively), see also refs. [40, 41] for previous bounds. The dashed blue line denotes a naive extrapolation of the expected bound from ref. [38] to the end of the high-luminosity LHC Run 5, assuming 3000 fb$^{-1}$ integrated luminosity at 14 TeV. For $c_{q}^{23} = c_{N}^{3}$ the two branching ratios of $W'$ are $\text{Br}(W' \to \tau \nu) : \text{Br}(W' \to 2j) \simeq 1 : 3$; the former is denoted by the horizontal grey dashed line in figure 4. The two branching ratios can be correspondingly smaller if other decay channels are open (for instance, to extra vector-like fermions, as contemplated in refs. [16, 17]). The grey shaded region is excluded by unitarity, which constrains $3(c_{q}^{23})^{2} + (c_{N}^{3})^{2} < 16\pi/g_{V}^{2}$ [42]. The experimental bounds shown in figure 4 assume that the $W'$ has a narrow width. This assumption fails for heavy $W'$ with a mass in the few TeV range. According to the results of a recast of the CMS search [39] performed for a wide $W'$ [43], the entire perturbative parameter space of the $W'$ model is excluded, except potentially for the very light $W'$, with masses below 500 GeV, where a reanalysis of older experiments would need to be carefully performed. Bounds on $W'$ from di-jet production [44–48] are less stringent and are not relevant for this simplified model.

Since the $W'_{\mu}$ couples to right-handed quarks, there is significant freedom in terms of the flavor structure of the $c_{q}^{ij}$ and $c_{N}^{i}$ couplings. We have limited the discussion to the
minimal case, taking only $c_{23}^q, c_N^3 \neq 0$, which is non-generic but possible, for instance, in flavorlocked models \cite{17,49}. In most flavor models all the $c_{ij}^q, c_N^i$ are non-zero, leading to constraints from precision measurements. In UV completions (see refs. \cite{16,17}), the $W'$ boson is expected to be accompanied by a $Z'$ state. The $Z'$ can, however, be parametrically heavier than the $W'$, in particular if additional sources of symmetry breaking are present. The collider constraints on $W'$ and $Z'$ are often comparable, while the flavor constraints from FCNCs are far more stringent for $Z'$ in the presence of any appreciable off-diagonal couplings \cite{17}: contributions from $W'$ exchange to flavor changing neutral currents only arise at one-loop and are significantly less constraining.

### 3.2 Vector leptoquark $U_1^\mu$

The interaction Lagrangian for the $U_1^\mu \sim (3,1)_{2/3}$ vector leptoquark is

$$\mathcal{L} \supset \alpha_{LQ}^{ij} \bar{L}^i_L \gamma_\mu Q^j_L U_1^{\mu 1} + \alpha_{\tilde{d}d}^{ij} \bar{\tilde{d}}^i_R \gamma_\mu d^j_R U_1^{\mu 1} + \alpha_{uN}^i \bar{u}_R^i \gamma_\mu N_R U_1^{\mu 1} + \text{h.c.},$$  \hspace{1cm} (3.4)

while the kinetic term, following the notation in \cite{50}, is

$$\mathcal{L} \supset -\frac{1}{2} U_{\mu\nu}^{\mu 1} U_{1\mu}^{\nu 1} + m_{U_1}^2 U_{1\mu}^{\mu 1} - ig_s \kappa U_{1\mu}^{\mu 1} T^a U_{1\nu} G^{a\mu\nu},$$  \hspace{1cm} (3.5)

with $U_{\mu\nu} = D_{\mu} U_{1\nu} - D_{\nu} U_{1\mu}$ the field strength tensor, and $\kappa$ a dimensionless coupling.

When the leptoquark is integrated out, eq. (3.4) gives two four-fermion operators, relevant for $R(D^{(*)})$ anomalies, with the Wilson coefficients

$$\frac{c_{SL}^{(\mu)}}{\rho_{SL} A_{eff}^2} = 2 \frac{\alpha_{LQ}^{33} \alpha_{uN}^{2}}{m_{U_1}^2}, \quad \frac{c_{VR}}{A_{eff}^2} = -\frac{\alpha_{\tilde{d}d}^{33} \alpha_{uN}^{2}}{m_{U_1}^2}. \hspace{1cm} (3.6)$$

The best fit values for the $U_1$ WCs in table 2 then imply

$$m_{U_1} \simeq 3.2 |\alpha_{LQ}^{33} \alpha_{uN}^{2}|^{1/2} \left[ \frac{40 \times 10^{-31}}{V_{cb}} \right]^{1/2} \text{TeV}, \hspace{1cm} (3.7)$$

with

$$\alpha_{\tilde{d}d}^{33} \simeq -5.8 \alpha_{LQ}^{33}. \hspace{1cm} (3.8)$$

where we used the lower set of best fits for $U_1$ in table 2 (the upper set is excluded by $B_c \rightarrow \tau \nu$, see figure 2). If one instead sets $c_{SL} = 0$, the best fit simply maps onto the $W'$ result (since both models then have the same non-zero coupling $c_{VR}$): $|c_{VR}| \simeq 0.46$, and

$$m_{U_1} \simeq 1.3 |\alpha_{\tilde{d}d}^{33} \alpha_{uN}^{2}|^{1/2} \left[ \frac{40 \times 10^{-31}}{V_{cb}} \right]^{1/2} \text{TeV}. \hspace{1cm} (3.9)$$

At the LHC, the $U_1$ leptoquark can be singly or pair produced. The pair production, $pp \rightarrow U_1 U_1^\dagger$, proceeds through gluon fusion, via the color octet term in (3.5), for which we take $\kappa = 1$ following ref. \cite{51}. The collider signatures of $U_1$ pair production depend on the
Figure 5. The LHC bounds from [51] (grey), [52] (brown), and [13, 53] (orange) on the $U_1^\mu$ vector leptoquark mass, assuming the relation $\alpha_{U_1}^{33} \simeq -5.8 \alpha_{LQ}^{33}$, arising from the $U_1$ best fit WCs to the $R(D^{(*)})$ data. Branching ratios for $U_1 \rightarrow c\nu, b\tau, t\nu$ decays are fixed by the remaining ratio of coupling constants $r_{U_1} = (\alpha_{uN}^2/\alpha_{LQ}^{33})^2$, assuming no other channels are open. Blue dashed lines denote contours satisfying the $U_1$ best fit mass relation (3.7) for $\alpha_{LQ}^{33} = 0.15, 0.3, 0.5, 1.0, \text{and} 2.0$.

$U_1$ decay channels. In the minimal set-up we switch on only three couplings, $\alpha_{LQ}^{33}, \alpha_{ld}^{33}$ and $\alpha_{uN}^2$, where $\alpha_{LQ}^{33}$ and $\alpha_{ld}^{33}$ are related through eq. (3.8), resulting in the branching ratios

$$
\text{Br}[U_1 \rightarrow t\bar{\nu}_\tau] : \text{Br}[U_1 \rightarrow b\tau] : \text{Br}[U_1 \rightarrow c\bar{N}_R] = |\alpha_{LQ}^{33}|^2 : (|\alpha_{LQ}^{33}|^2 + |\alpha_{ld}^{33}|^2) : |\alpha_{uN}^2|^2 \quad (3.10)
$$

$$
= \frac{0.03}{1 + 0.03 r_{U_1}} : \frac{0.97}{1 + 0.03 r_{U_1}} : \frac{0.03 r_{U_1}}{1 + 0.03 r_{U_1}},
$$

where

$$
r_{U_1} = \left(\frac{\alpha_{uN}^2}{\alpha_{LQ}^{33}}\right)^2.
$$

Here, for simplicity, we have neglected the final state masses and the small corrections due to the off-diagonal CKM matrix elements in the $\alpha_{LQ}^{ij}(E_{L}^i, u_{Q_j}^i, U_y^U_1)$. The presence of left-handed quark doublets also inevitably leads to CKM suppressed transitions $U_1 \rightarrow c\bar{\nu}_\tau, u\bar{\nu}_\tau, s\tau, d\tau$.

The corresponding LHC bounds for $U_1$ are shown in figure 5, assuming no other decay channels are open. The most stringent bounds come from $pp \rightarrow U_1 U_1$ pair production, with both leptoquarks decaying either as $U_1 \rightarrow c\bar{N}_R$ [51] (grey region) or $U_1 \rightarrow b\tau$ [52] (brown region). Ref. [51] also gives bounds for the decay channel $U_1 \rightarrow t\nu_\tau$, which are not shown in figure 5 as they are always weaker in our setup. We see that direct searches still allow for $m_{U_1} \geq 1.5 \text{ TeV}$, where the parameters of the model are still perturbative, as an explanation for the $R(D^{(*)})$ anomalies. It is worth noting that a simultaneous fit to all three decay channels by the experiments would improve the sensitivity to $U_1$; such an analysis is likely the most optimal strategy for discovering a $U_1$ state responsible for the $R(D^{(*)})$ anomalies.

Figure 5 also shows the constraint on the $U_1$ model parameter space from the CMS $pp \rightarrow \tau\tau$ search [53] (see also ATLAS search [54]). In orange is shown the constraint on $r_{U_1}$, as a function of $m_{U_1}$, that is obtained from figure 6 of ref. [13] with the replacement
$g_U \rightarrow \left[(\alpha_{LQ}^{33})^2 + (\alpha_{d}^{3d})^2\right]^{1/2}$. Assuming the relation $\alpha_{LQ}^{33} \simeq -5.8 \alpha_{d}^{3d}$, arising from the $U_1$ best fit WCs to the $R(D^*)$ data, the bound on $g_U$ in [13] translates to the excluded region in figure 5.

3.3 Scalar leptoquark $S_1$

The scalar leptoquark $S_1 \sim (3, 1)_{1/3}$ has the following interaction Lagrangian,

$$\mathcal{L} \supset z_u(U_R^T \ell R)S_1 + z_d(\bar{d}_R N_R)S_1 + z_Q(Q_L^T \ell L_L)S_1.$$  \hspace{1cm} (3.12)

Integrating out the leptoquark generates the following interaction Lagrangian above the electroweak scale

$$\mathcal{L}_{\text{eff}}^S = \frac{z_d z_u^*}{2 m_S^2} Q_{\text{VR}} - \frac{z_d z_d^*}{2 m_S^2} \left(Q_{\text{SR}} - \frac{1}{4} Q_T\right) + \frac{z_u z_u^*}{2 m_S^2} \left[\epsilon_{ab}(\bar{R}_L^T u_R Q_b^T - \frac{1}{4} \epsilon_{ab}(\bar{R}_L^T \sigma_{\mu\nu} L_L^T)(\bar{u}_R \sigma_{\mu\nu} Q_b^T)\right] + \text{h.c.},$$  \hspace{1cm} (3.13)

where the operators $Q_{\text{VR}}, Q_{\text{SR}}, Q_T$ are defined in (2.3). The $b \rightarrow c\tau \bar{N}_R$ decay is generated if $z_u^2 z_d^2 \neq 0$ or $z_Q^2 z_d^2 \neq 0$. The two operators in the second line give rise to the $b \rightarrow c\tau \nu_i$ decay for $z_Q^2 z_u^2 \neq 0$, where $\nu_i$ are the SM neutrinos, which interfere with the SM contribution; for simplicity, we therefore only consider the $b \rightarrow c\tau \bar{N}_R$ decay, setting $z_Q^2 = 0$, so that only the operators in the first line in (3.13) are generated (alternatively, one may consider the regime $z_u, z_Q \ll z_d$, so that the contribution from the second line is negligible).

In the analysis of collider constraints, we conservatively keep only the minimal set of $S_1$ couplings required for the $R(D^*)$ anomaly nonzero: $z_u^2 z_d^2, z_Q^2 z_d^2 \neq 0$. The Wilson coefficients of the $b \rightarrow c\tau \bar{N}_R$ operators $\mathcal{O}_{\text{VR}}, \mathcal{O}_{\text{SR}}, \mathcal{O}_T$ are given by,

$$\frac{c_{\text{VR}}}{\Lambda_{\text{eff}}^2} = -\frac{z_u^2 z_d^2}{2 m_S^2}, \quad \frac{c_{\text{SR}}}{\Lambda_{\text{eff}}^2} = -\frac{z_Q^2 z_d^2}{2 m_S^2}.$$  \hspace{1cm} (3.14)

The best fit values for the $S_1$ WCs in table 2 then imply

$$m_{S_1} \approx 1.2 |z_u^2 z_d^2|^{1/2} \left[\frac{40 \times 10^{-3}}{V_{cb}}\right]^{1/2} \text{TeV},$$  \hspace{1cm} (3.15)

with

$$z_u^2 \approx 1.1 z_Q^2.$$  \hspace{1cm} (3.16)

using the lower set of best fits for $S_1$ in table 2 (the upper set is excluded by $B_c \rightarrow \tau \nu$, see figure 2). The branching ratios for $S_1$ decays are thus

$$\text{Br}[S_1 \rightarrow c\tau] : \text{Br}[S_1 \rightarrow bN_R] : \text{Br}[S_1 \rightarrow s\nu_\tau] = \left( |z_u^2|^2 + |z_Q^2| \right) : |z_d^2| : |z_Q^2|$$

$$\approx 0.69 : 0.37 r_{S_1} : 0.31$$

$$\frac{1}{1 + 0.37 r_{S_1}} : \frac{1}{1 + 0.37 r_{S_1}} : \frac{1}{1 + 0.37 r_{S_1}},$$  \hspace{1cm} (3.17)
where we have defined

\[ r_{S_1} = \left( \frac{z_2^3}{z_u} \right)^2. \]  

(3.18)

The resulting bounds from \( pp \to S_1S_1 \) pair production at the 13 TeV LHC are shown in figure 6. The grey shaded region is excluded by the CMS search [51] with 35.9 fb\(^{-1}\) integrated luminosity, assuming both \( S_1 \) decay as \( S_1 \to bN_R \) with the branching ratio in (3.17). The brown shaded region is excluded by the CMS search [52] using 12.9 fb\(^{-1}\) integrated luminosity, assuming \( pp \to S_1S_1 \) followed by \( S_1 \to c\tau \) decay, with the \( r_{du} \) dependent branching ratio in (3.17). We have assumed the \( S_1 \) best fit mass relation (3.16) to \( R(D^{(*)}) \) data to derive these bounds.

The orange shaded region in figure 6 shows the 95% CL constraint from the recast of the 13 TeV ATLAS \( pp \to \tau\tau \) search at 36 fb\(^{-1}\) integrated luminosity [54], performed in ref. [55]. The bounds in figure 3 (left) in ref. [55] can be reinterpreted in terms of the \( S_1 \) model coupling to a right-handed neutrino by making the replacement \( \lambda_{21}^2 \to \left[ (z_2^3)^2 + (z_Q^2)^2 \right]^{1/2} \).

The combined set of constraints indicates that the \( S_1 \) leptoquark can be consistent with the \( R(D^{(*)}) \) anomaly for \( m_{S_1} \) as low as 1000 GeV, and with perturbative couplings (the required values of \( z_u^3 \) are shown by dashed blue lines in figure 6).

### 3.4 Scalar leptoquark \( \tilde{R}_2 \)

The scalar leptoquark \( \tilde{R}_2 \sim (3, 2)_{1/6} \) has the following interaction Lagrangian,

\[ \mathcal{L} \supset \alpha_L d \left( \bar{L}_L d_R \right) \tilde{R}_2^3 + \alpha_Q (\bar{Q}_L N_R) \tilde{R}_2 + \text{h.c.}. \]  

(3.19)

Integrating out the \( \tilde{R}_2 \) generates

\[ \frac{c_{SR}^{(\mu)}}{\rho_{SR} \lambda_{eff}^2} = \frac{c_T^{(\mu)}}{\rho_T \lambda_{eff}^2} = \frac{\alpha_{Ld} \alpha_{QN}}{2 m_{\tilde{R}_2}^2}. \]  

(3.20)
In this section, we discuss the phenomenology associated with the right-handed (sterile) neutrino $N_R$. As we will see below, the coupling of $N_R$ to the SM fermions through one of
the higher dimension operators in eq. (2.8), needed to explain $R(D^{(*)})$, carries interesting implications for neutrino masses, cosmology, and collider signatures. We will assume that $N_R$ is a Majorana fermion with mass $\lesssim \mathcal{O}(100)$ MeV so that it remains compatible with the measured missing invariant mass spectrum in the $B \to D^{(*)}\tau\nu$ decay chain. As in section 3, we do not consider the $\Phi$ model as it is excluded by $B_c \to \tau\nu$ constraints.

4.1 Neutrino masses

The effective operators (2.8) induce a $N_R - \nu_L$ Dirac mass at the two loop order via contributions of the form

$$m_D N_R \nu_L \sim W \cdot \begin{array}{c}
\tau \\
N_R \\
\nu_L \\
\bar{b} \\
W \\
\tau
\end{array}.$$  

Here, the simplified model mediator has been integrated out, producing an effective four-fermion vertex, shown in gray. Depending on the chiral structure of the simplified model, various mass insertions are mandated on the internal quark and lepton lines. In particular, the $O_{\text{VR}}$ operator requires three mass insertions, while the scalar and tensor operators require only one. The corresponding Dirac masses can be estimated as

$$W': \quad m_D \sim \frac{c_{\text{VR}} g_2^2}{2 \Lambda_{\text{eff}}^2} \frac{V_{cb}}{(16\pi^2)^2} m_b m_c m_\tau \sim c_{\text{VR}} 10^{-3} \text{eV},$$

$$R_2: \quad m_D \sim c_{SR} m_b g_2^2 \frac{V_{cb}}{(16\pi^2)^2} \sim c_{SR} 10^2 \text{eV},$$

$$U_1: \quad m_D \sim \left[ c_{\text{SL}} m_c + \frac{c_{\text{VR}}}{\Lambda_{\text{eff}}^2} m_b m_c m_\tau \right] \frac{g_2^2}{2} \frac{V_{cb}}{(16\pi^2)^2} \sim (c_{\text{SL}} 10^2 + c_{\text{VR}} 10^{-3}) eV,$$

$$S_1: \quad m_D \sim \left[ c_{\text{SR}} m_b + \frac{c_{\text{VR}}}{\Lambda_{\text{eff}}^2} m_b m_c m_\tau \right] \frac{g_2^2}{2} \frac{V_{cb}}{(16\pi^2)^2} \sim (c_{\text{SR}} 10^2 + c_{\text{VR}} 10^{-3}) eV.$$  

In the above estimates, we have ignored $O(1)$ prefactors and loop integral factors apart from those implied by naive dimensional analysis. Note that for diagrams with a single mass insertion, the Wilson coefficients $c_{\text{SL}}$, $c_{\text{SR}}$ appear without the $1/\Lambda_{\text{eff}}^2$ prefactor. In such cases, strictly speaking, it is the couplings of the mediators rather than the Wilson coefficients that should appear in the estimates. However, since the collider constraints require mediators to be heavy, with mass approximately equal to $\Lambda_{\text{eff}}$, it is a reasonable approximation to use the Wilson coefficients everywhere in the above estimates.

Furthermore, for $R_2$, $U_1$, and $S_1$ mediators, which couple to the left-handed $\tau_L$, there are additional two loop contributions to the neutrino mass matrix arising from the $SU(2)_L$ related operators involving $\nu_L$. A representative diagram is shown in figure 8. While such diagrams contain similar mass insertions and WC scalings as the corresponding $c_{\text{SL,SR}}$ terms in eqs. (4.2), they are GIM suppressed and thus expected to produce only subleading corrections to the Dirac mass estimates in eqs. (4.2).
Since $N_R$ is assumed to have a Majorana mass $m_{N_R} \lesssim 100$ MeV, the contribution to the SM neutrino masses is $\sim m_D^2/m_{N_R}$, which should not exceed the observed neutrino mass scale $m_\nu \sim 0.1$ eV. From the best fit regions shown in figures 2 or 3 (and the best fit values from table 2), it follows that the $W'$-mediated diagram gives a Dirac mass $m_D \sim 10^{-3}$ eV, which is consistent with observed neutrino masses, whereas the $R_2$ mediated diagram gives $m_D \sim 100$ eV, which is in some tension for $m_{N_R} \lesssim 10$ keV. Likewise, the $U_1$ and $S_1$ models produce similarly problematic contributions to the neutrino masses at their best fit points (see table 2). However, from figures 2 and 3 we also see that the $1\sigma$ CLs of the $U_1$ and $S_1$ models do contain regions with the scalar Wilson coefficients $|c_{SL,SR}| \ll 1$, corresponding to small couplings $\alpha_{LQ} \ll 1$ and $z_Q \ll 1$ (cf. eqs. (3.6) and (3.14)), which remain compatible with observed neutrino masses.

If additional operators are present, neutrino mass contributions can also be generated at one loop. For instance, as discussed in section 2.2, new operators coupling to second generation quark doublets can be introduced to cancel away large contributions to $b \to s\nu \bar{\nu}$ from the operators in eq. (2.11). Such 1-loop neutrino mass contributions scale as $m \sim \frac{1}{16\pi^2} m_f$ and, depending on whether the new operators couple to $\nu\nu$ or $\nu N_R$, contribute to the Majorana or Dirac mass terms for the neutrinos. Unless suppressed by small couplings in the diagram, such mass contributions are generally several orders of magnitude larger than what is allowed by the observed neutrino mass scale $m_\nu \sim 0.1$ eV, and would need to be cancelled by fine-tuned values of bare neutrino masses.

Additional Dirac mass contributions beyond the diagrams considered above could worsen or improve the outlook. For instance, if the mediators also couple to other quarks, in particular the top quark, the corresponding two loop diagrams with a top quark mass insertion would lead to unacceptably large contributions to neutrino masses. On the other hand, additional Dirac mass terms that interfere destructively with the two loop contributions here could restore consistency in otherwise problematic regions of parameter space, albeit at the cost of some fine-tuning of parameters.

### 4.2 Sterile neutrino decay

The two loop diagrams considered above also give rise to the decay process $N_R \to \nu \gamma$ via the emission of a photon from one of the internal propagator lines (a representative diagram...
Figure 9. Sterile neutrino decay modes induced by the NP couplings (left) and by tree level sterile-active mixing (centre, right).

\[ \begin{array}{|c|c|c|}
\hline
\text{Model} & \Gamma_{N_R \rightarrow \nu \gamma} & \text{lifetime (s)} \\
\hline
W' & \frac{\alpha^2_{\text{VR}}}{\Lambda^2_{\text{eff}}} \alpha' \gamma^{2} G_F^2 m^2_\nu m^2_b m^2_\ell m^3_{N_R} & c^2_{\text{VR}} \frac{10^{24}}{(m_{N_R}/\text{keV})^{-3}} \\
\hat{R}_2 & \frac{\alpha_{\text{SR}}}{\Lambda^2_{\text{eff}}} \alpha' \gamma^{2} G_F^2 m^2_\nu m^3_{N_R} & c^2_{\text{SR}} \frac{10^{13}}{(m_{N_R}/\text{keV})^{-3}} \\
U_1 & \frac{\alpha_{\text{SL}}}{\Lambda^2_{\text{eff}}} \alpha' \gamma^{2} G_F^2 m^2_\nu m^3_{N_R} & c^2_{\text{SL}} \frac{10^{14}}{(m_{N_R}/\text{keV})^{-3}} \\
S_1 & \frac{\alpha_{\text{SR}}}{\Lambda^2_{\text{eff}}} \alpha' \gamma^{2} G_F^2 m^2_\nu m^3_{N_R} & c^2_{\text{SR}} \frac{10^{13}}{(m_{N_R}/\text{keV})^{-3}} \\
\hline
\end{array} \]

Table 3. Approximate $N_R \rightarrow \nu \gamma$ decay rates (middle column) and lifetimes (final column) for the mediators listed in the first column. For $U_1(S_1)$, we only show the contribution from the $c_{\text{SL}}(c_{\text{SR}})$ operators, which are expected to dominate; if these coefficients vanish, the decay rates and lifetimes get contributions from $c_{\text{VR}}$ of the same form as that for the $W'$ operator.

is shown in figure 9 (left)). The approximate $N_R \rightarrow \nu \gamma$ decay rates\(^1\) for the simplified models, along with the corresponding decay lifetime estimates, are listed in table 3 (for related calculations, see refs. [56–59]). Note that for a given mediator and sterile neutrino mass $m_{N_R}$, the decay rate is completely fixed by the Wilson coefficients consistent with the $R(D^{(*)})$ anomaly.

For appreciable mixing between $N_R$ and the SM neutrinos, the leading tree-level decay is into three SM neutrinos (figure 9 center) and, if kinematically accessible, into charged leptons (figure 9 right). The $N_R \rightarrow 3\nu$ decay rate is

\[ \Gamma_{N_R \rightarrow 3\nu} \approx \frac{G_F^2}{192 \pi^3} m_{N_R}^5 \sin^2 \theta \approx 10^{-48} \left( \frac{m_{N_R}}{\text{keV}} \right)^5 \left( \frac{\sin^2 \theta}{10^{-4}} \right) \text{GeV}, \]  

(4.3)

where $\theta$ is the mixing angle between $N_R$ and the SM neutrino. The $N_R \rightarrow 3\nu$ decay width is in general subdominant to the $N_R \rightarrow \nu \gamma$ decay width induced by the $R(D^{(*)})$ anomaly. For a direct comparison, one can rewrite the $N_R \rightarrow \nu \gamma$ decay rate in table 3 in terms of the Dirac mass from eq. 4.2, then convert to the mixing angle via $\sin \theta \approx m_D/m_N$. For instance, for $S_1$ this gives $\Gamma(N \rightarrow \nu \gamma) \approx 32 \alpha \sin^2 \theta m^3_N G_F^2 / \pi^4 / g^4$. Thus

\[ \frac{\Gamma(N \rightarrow \nu \gamma)_{S_1}}{\Gamma(N_R \rightarrow 3\nu)_{S_1}} \approx \frac{32 \times 192 \alpha}{\pi g^4} \approx 10^3. \]  

(4.4)

\(^1\)The mass insertion required by the helicity flip for the emission of a photon can occur on an internal fermion line, and does not incur the cost of a mass suppression on an external fermion leg, in contrast to $f_1 \rightarrow f_2 \gamma$ diagrams via an SU(2)$_L$ electroweak loop.
4.3 Sterile neutrino cosmology

The above estimates imply that the sterile neutrino $N_R$ can be fairly long-lived. The interactions with SM fermions mandated by consistency with the $R(D^{(*)})$ anomaly also lead to copious production of $N_R$ in the early Universe. The cosmological aspects of the sterile neutrino therefore require careful treatment.

The interactions with SM fermions thermalize the $N_R$ population with the SM bath at high temperatures. These interactions are active until the temperature drops below the masses of the SM fermions involved in these interactions, i.e., around the GeV scale. Since we have assumed $m_{N_R} \lesssim 100$ MeV, the $N_R$ abundance is not Boltzmann suppressed, and $N_R$ survives as an additional relativistic neutrino species in the early Universe. It then becomes crucial to determine the fate of this $N_R$ population.

For the $R_2$, $U_1$, and $S_1$ mediated models, it follows from table 3 that the $N_R$ lifetime is $\sim 10^{14}(m_{N_R}/\text{keV})^{-3}$ s. For $m_{N_R} \sim \mathcal{O}(\text{eV–keV})$, this implies a late decay of the $N_R$ population into the $\gamma\nu$ channel, which injects an unacceptable amount of photons into the diffuse photon background. The exception are masses close to the upper limit of the range we consider, $m_{N_R} \lesssim 100$ MeV, for which the lifetime is reduced to $\lesssim 1$ s. The decays then occur before big bang nucleosynthesis (BBN) and do not leave any visible imprints.

In contrast, for the $W'$ mediated case (or for $U_1$, $S_1$ in the parts of the Wilson coefficient $1\sigma$ CL regions where $c_{\text{SL}}, c_{\text{SR}}$ are vanishingly small), the lifetime is much longer because of the additional mass insertions in the decay diagrams, and a lifetime $\lesssim 1$ s cannot be achieved for any realistic choices of parameters. However, for $m_{N_R} \lesssim 100$ keV, the sterile neutrino has a lifetime greater than the age of the Universe and could in principle form a component of dark matter or dark radiation.

The dark matter and dark radiation possibilities of $N_R$ in the $W'$ model have been extensively discussed in ref. [17]. In contrast to traditionally studied frameworks of sterile neutrino dark matter, where the relic abundance is produced via freeze-in mechanisms (see, e.g., [60–65]), the $W'$ model involves the sterile neutrino freezing out as a relativistic species, leading to too large of a relic abundance for masses greater than $\mathcal{O}(\text{keV})$. This can be fixed with appropriate entropy dilution from, for instance, late decays of GeV scale sterile neutrinos [58, 58, 66, 67], which also makes the dark matter colder, improving compatibility with warm dark matter constraints. The $\gamma$-ray bounds from various observations [68] rule out dark matter lifetimes of $\mathcal{O}(10^{26–28})$ s in the keV–MeV window, ruling out the case that $N_R$ constitutes all of dark matter. This leaves us with the possibility that $N_R$ may constitutes a small fraction — at the sub-percent level — of dark matter. Future $\gamma$-ray observations will probe this possibility and could discover a line signal from the $N_R \to \gamma\nu$ decay. For masses $m_{N_R} \lesssim \text{keV}$, $N_R$ can act as dark radiation and contribute to the effective number of relativistic degrees of freedom $\Delta N_{\text{eff}} \approx \mathcal{O}(0.1)$ at BBN and/or CMB decoupling, which could be detected with future instruments such as CMB-S4 [69]. Lifetimes shorter than the age of the Universe, however, are incompatible with current observational constraints.
4.4 Displaced decays at direct searches and colliders

As discussed in the previous section, in the $R_2, U_1,$ and $S_1$ models, cosmology favors the regime $m_{N_R} \sim 100$ MeV, with a lifetime $\lesssim 1\,\text{s}$. Since the dominant decay channel is $N_R \to \nu \gamma$, this would give rise to displaced decays into a photon+MET. Such displaced signals could provide an interesting, but challenging, target for proposed detectors such as SHiP [70], MATHUSLA [71], FASER [72], and CODEX-b [73]. Displaced decays can also occur in the $W'$ UV completion of refs. [16, 17], where, as discussed earlier, GeV scale sterile neutrinos with lifetimes $\lesssim 1\,\text{s}$ might be needed to entropy dilute problematic overabundances of the $N_R$; these can also lead to several other observable signals at various direct and cosmological probes (see, e.g., the discussion in [74]).

5 Conclusions

We have performed an EFT study of the lowest dimension electroweak operators that can account for the $R(D^{(*)})$ anomalies, assuming they arise because of incoherent contributions from semitauonic decays involving a right-handed sterile neutrino $N_R$. These dimension-six operators can arise from a tree-level mediator exchange in five possible simplified models. We examined the fits and constraints for each simplified model. While all five models have $1\sigma$ fit regions consistent with the $R(D^{(*)})$ data, the case of the scalar doublet mediator is conservatively in tension with constraints from $\text{Br}[B_c \to \tau \nu]$, while the experimental bounds on $b \to s \nu \bar{\nu}$ rates are in tension with the predicted rates from the scalar leptoquark $R_2$.

The fit regions of the remaining three simplified models imply sizable semileptonic branching ratios for the tree-level mediators. We find that each model already faces fairly stringent collider constraints. The searches for the $W'$ mediator in the $W' \to \tau \nu$ channel exclude the model for perturbative couplings, where the calculations are reliable, with the possible exception of very light $W'$ masses (see figure 4 and surrounding discussion). The two leptoquark models are consistent with LHC search results provided the mediator masses are $\mathcal{O}(\text{TeV})$, while their couplings may still remain in the perturbative regime. Our analysis indicates promising paths to future discovery of the tree-level mediators at the LHC, with couplings and masses consistent with the fit to the $R(D^{(*)})$ data. The vector leptoquark $U_1'$ can best be probed at the LHC with simultaneous fits to the three decays $U_1 \to cN_R, U_1 \to b\tau$ and $U_1 \to t\nu_\tau$. Likewise, the scalar leptoquark $S_1$ can be probed via $S_1 \to bN_R$ and $S_1 \to c\tau$ decays. Since the mediators cannot be arbitrarily heavy if the couplings are to remain perturbative, prospects of detecting them at the LHC are quite encouraging.

We have also discussed the phenomenology associated with the sterile neutrino $N_R$. In simplified models involving $R_2, U_1,$ and $S_1$, constraints from contributions to neutrino masses as well as cosmology indicate a preference for $m_{N_R} \sim 10$–$100$ MeV with a decay lifetime $\lesssim 1\,\text{s}$ in the dominant channel $N_R \to \nu \gamma$. This opens up the potential for detecting displaced decays of $N_R$ at various detectors. It also implies potentially measurable distortions of the kinematical distributions in semileptonic $B$ meson decays due to the heavy sterile neutrino in the final state. For the $W'$ simplified model, the predicted contribution to neutrino masses is much smaller and poses no constraints on the model. The
predicted decay lifetime of $N_R$ is correspondingly much longer than the age of the Universe. Consequently, a significant relic abundance of $N_R$ is likely present in the universe, which can contribute to dark radiation and give measurable deviations to the effective number of relativistic degrees of freedom $\Delta N_{\text{eff}} \approx O(0.1)$ at BBN and/or CMB decoupling for $m_{N_R} \lesssim \text{keV}$, or constitute a small fraction of dark matter for $N_R$ in the keV–MeV mass range with possible gamma ray signals at future probes.

The interpretation of the $R(D^{(*)})$ anomaly in terms of new physics coupling the SM fermions to a right-handed sterile neutrino is therefore an exciting possibility with testable predictions in multiple directions, spanning kinematic distributions of the measured $B$ meson decays, searches for heavy TeV scale particles at the LHC, displaced decay signals at various detectors, as well as astrophysical and cosmological signatures.

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A Differential distributions

In this appendix we collect the predictions for several normalized differential distributions for $B \to (D^{(*)} \to D\pi)(\tau \to \ell\nu\bar{\nu})\bar{\nu}$ and $B \to D(\tau \to \ell\nu\bar{\nu})\bar{\nu}$ decay chains, shown in the left and right columns in figures 10–13, respectively. In each plot, the SM predictions (blue dashed curves) are compared with the predictions for the particular simplified model (grey bands), obtained by varying the relevant Wilson coefficients over the $2\sigma$ regions in figure 2. In each of the figures the first row shows the normalized distribution $(1/\Gamma)(d\Gamma/dE_D)$, where $E_D$ is the energy of the outgoing $D$ meson in the $B$ meson rest frame. The second row contains the $(1/\Gamma)(d\Gamma/dE_\ell)$ distribution, with $E_\ell$ the energy of the final state charged lepton, while the third row shows the $(1/\Gamma)(d\Gamma/dm_{\text{miss}}^2)$ distribution, with $m_{\text{miss}}^2$ the combined invariant mass of the system of three final state neutrinos. The final row in each figure shows the $(1/\Gamma)(d\Gamma/d\cos \theta_{D\ell})$ normalized distribution, where $\theta_{D\ell}$ is the angle between the three momenta of the $D$ meson and the charged lepton, $\ell$, in the rest frame of the $B$ meson.

The comparison between the SM predictions (blue dashed curves) and the predictions for the $W'$ simplified model (grey bands) is shown in figure 10. The differences between the two predictions are small, below about 10% for $(1/\Gamma)(d\Gamma/dE_\ell)$ and well below this for the other distributions. Similarly small corrections from NP to the shapes of distributions are found for the $R_2$ model, figure 11. In this case the largest deviation is found for the $(1/\Gamma)(d\Gamma/dE_D)$ distribution for the $\bar{B} \to D^* \to D\pi$ decay (figure 11, first row, right panel).
Figure 10. Gray bands show kinematic distributions for $B \rightarrow (D^* \rightarrow D\pi)(\tau \rightarrow \ell\nu\ell\nu)\bar{\nu}^{}$ (left) and $B \rightarrow D(\tau \rightarrow \ell\nu\ell\nu)\bar{\nu}^{}$ (right) in the $B$ rest frame for the $W'$ simplified model in table 1, with the Wilson coefficient $c_{VR}^{}$ ranging over $2\sigma$ best fit regions in figure 2, and applying the phase space cuts (2.20). The blue dashed curves show the SM prediction.

and is at the level of about $\mathcal{O}(20\%)$. The deviations are potentially sizable for the $U_1$ and $S_1$ models for at least some of the distributions, see figures 12 and 13, respectively.
Figure 11. Gray bands show kinematic distributions for $B \to (D^* \to D\pi)(\tau \to \ell\bar{\nu}_\ell\nu_\tau)\bar{\nu}$ (left) and $B \to D(\tau \to \ell\bar{\nu}_\ell\nu_\tau)\bar{\nu}$ (right) in the $B$ rest frame for the $\tilde{R}_2$ simplified model in table 1, with the Wilson coefficients $c_{SR} = 4c_T$ ranging over $2\sigma$ best fit regions in figure 2, and applying the phase space cuts (2.20). The blue dashed curves show the SM prediction.
Figure 12. Gray bands show kinematic distributions for $\bar{B} \rightarrow (D^* \rightarrow D\pi)(\tau \rightarrow \ell\bar{\nu}_\ell\nu_\tau)\bar{\nu}$ (left) and $\bar{B} \rightarrow D(\tau \rightarrow \ell\bar{\nu}_\ell\nu_\tau)\bar{\nu}$ (right) in the $B$ rest frame for the $U_1$ simplified model in table 1, with the Wilson coefficients $c_{SL}, c_{VR}$ ranging over $2\sigma$ best fit regions in figure 2, and applying the phase space cuts (2.20). The blue dashed curves show the SM prediction.
Figure 13. Gray bands show kinematic distributions for $B \to (D^* \to D \pi)(\tau \to \ell \bar{\nu}_\ell \nu_{\tau})\bar{\nu}$ (left) and $B \to D(\tau \to \ell \bar{\nu}_\ell \nu_{\tau})\bar{\nu}$ (right) in the $B$ rest frame for the $S_1$ simplified model in table 1, with the Wilson coefficients $c_{\text{VR}}^2, c_{\text{SR}}^2 = 4c_T$ ranging over 2σ best fit regions in figure 2, and applying the phase space cuts (2.20). The blue dashed curves show the SM prediction.
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