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$b \to c\tau\bar{\nu}_{e,\mu}$ contributions to $R(D^{(*)})$

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Abstract: The $R(D^{(*)})$ puzzle stands for a $\sim 3\sigma$ violation of lepton flavor universality between the decay rates of $B \to D^{(*)}\tau\nu$ and $B \to D^{(*)}\ell\nu$, where $\ell = e, \mu$. If it is accounted for by new physics, there is no reason in general that the relevant neutrinos are, respectively, ν_τ and ν_ℓ. We study whether the τ related rate could be enhanced by significant contributions to $B \to D^{(*)}\tau\nu_\ell$ from a class of operators in the Standard Model Effective Field Theory (SMEFT). We find the upper bounds from forbidden or rare meson decays imply that the contributions from the lepton flavor violating processes account for no more than about 4% of the required shift. Yet, no fine-tuned flavor alignment is required for the new physics. Searching for the related high-p_T process $pp \to \tau^\pm\mu^\mp$ can at present put a lower bound on the scale of the lepton flavor violating new physics that is a factor of 2.2 weaker than the bound from meson decays. An exception to our conclusion arises from a specific combination of scalar and tensor SMEFT operators.

Keywords: Semi-Leptonic Decays, SMEFT, Specific BSM Phenomenology

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JHEP07(2022)077

Contents

JHEP07(2022)077

1 Introduction

Within the Standard Model (SM), the electroweak interactions of the leptons are flavor universal. Violation of lepton flavor universality arises from Yukawa interactions, that are negligible in this context, and from phase space effects, which are calculable. A test of the SM prediction of lepton flavor universality is provided by the ratios

$$R(D^{(*)}) \equiv \frac{\Gamma(B \to D^{(*)}\tau\bar{\nu})}{\Gamma(B \to D^{(*)}\ell\bar{\nu})}, \quad (\ell = e \text{ or } \mu). \tag{1.1}$$

The SM predictions, derived by naive averaging [1] over the results reported in refs. [2–5], are

$$\begin{aligned} R(D) &= 0.299 \pm 0.003, \\ R(D^*) &= 0.258 \pm 0.005. \end{aligned} \tag{1.2}$$

The current world averages for $R(D)$ and $R(D^*)$, combining the results reported in refs. [6–14] are as follows [1]:

$$\begin{aligned} R(D) &= 0.340 \pm 0.030, \\ R(D^*) &= 0.295 \pm 0.014. \end{aligned} \tag{1.3}$$

The difference of the experimental measurements from the SM predictions corresponds to about 3.1σ (p-value of 2.7×10^{-3}). We thus aim to explain

$$R(D^{(*)})/R(D^{(*)})^{\text{SM}} \approx 1.14 \pm 0.05. \tag{1.4}$$

The quark transition via which the $B \to D^{(*)}\tau\nu$ proceeds is $b \to c\tau\nu$. Note, however, that the flavor of the neutrino is, of course, unobservable. It could be ν_τ, in which case

the process respects the accidental lepton flavor symmetry of the SM. There is no reason, however, that the symmetry is respected by new physics, particularly when the new physics violates lepton flavor universality, so that the neutrino could also be ν_μ or ν_e or some combination of the three flavors. This possibility has been discussed very little in the literature (for an exception, see [15]), and we aim to fill in this gap. We ask three main questions:

- Could the $R(D^{(*)})$ puzzle be solved via new physics contributions to $b \to c\tau\nu_\ell$ with $\ell = e, \mu$?
- If not, how precise should the alignment of ν with ν_τ be?
- What is the sensitivity of the high-p_T experiments at the LHC to such lepton flavor violating new physics?

The plan of this paper is as follows. In section 2, we present the theoretical framework of the SM effective field theory (SMEFT), within which we carry out our analysis, and estimate the size of the dimension-six operators that can explain the $R(D^*)$ puzzle. In section 3, we obtain bounds on these operators from various processes to which they contribute. In section 4 we explore the reach of current and future collider experiments to probe the dimension-six terms with searches of lepton flavor violating di-lepton final states. We summarize our conclusions in section 5. A discussion of additional operators is given in appendix A.

2 $R(D^{(*)})$ in the SMEFT

We assume that the new physics contributions originate at a scale $\Lambda \gg v$, and consider the following two terms in the SMEFT Lagrangian [15]:

$$\mathcal{L}_{\rm NP} = \frac{C_1^{ilkm}}{\Lambda^2}(\overline{L_i}\gamma_\sigma L_l)(\overline{Q_k}\gamma^\sigma Q_m) + \frac{C_3^{ilkm}}{\Lambda^2}(\overline{L_i}\gamma_\sigma \tau^a L_l)(\overline{Q_k}\gamma^\sigma \tau^a Q_m), \tag{2.1}$$

where L is the SU(2)-doublet lepton field, Q is the SU(2)-doublet quark field, and i, l, k, m are flavor indices. For the sake of definiteness, and to avoid the strongest constraints from flavor changing neutral current (FCNC) processes, we take $i = \tau$, $k = s$, and $m = b$, while l runs over e, μ, τ. Three comments are in order concerning our choices for the flavor, Lorentz and CP structures:

- The weakest constraints apply when $k = m = b$, in which case no FCNC in the down sector are generated. However, the contribution to the $b \to c\tau\bar{\nu}_\ell$ decay rate gets an extra suppression by $|V_{cb}|^2$ compared to the $k = s$ case. This brings the relevant new physics scale close to the electroweak scale (roughly, $(0.14)^{-1/4} m_W$), a situation that does not lend itself to an SMEFT analysis and requires a model dependent analysis instead (see, e.g. ref. [16] for relevant simplified models). Bounds from Υ decays [17–19] are relevant for this scenario.

JHEP07(2022)077

- The contributions from individual scalar or tensor operators are disfavored by their different contributions to $R(D)$ and $R(D^*)$, by the B_c lifetime and by their modification of the differential decay rates with respect to the SM [20–23]. An exception to this statement is provided by a specific combination of scalar and tensor operators. We present further details on this issue in appendix A.

- Given that we focus on lepton flavor violating operators, there is no interference with the SM operators. Thus, the measurements that we discuss are sensitive only to absolute values of Wilson coefficients, and not to their phase structure.

We denote $C_{1,3}^{\tau lsb}$ by $C_{1,3}^{l}$. The $C_{1,3}^{l}$-dependent terms can be rewritten as follows:

$$\begin{aligned}\Lambda^2\mathcal{L}_{\rm NP} = &(C_1^l + C_3^l)V_{is}V_{jb}^*(\overline{u_{Li}}\gamma^\mu u_{Lj})(\overline{\nu_\tau}\gamma_\mu\nu_l)\\ &+(C_1^l - C_3^l)V_{is}V_{jb}^*(\overline{u_{Li}}\gamma^\mu u_{Lj})(\overline{\tau_L}\gamma_\mu l_L)\\ &+(C_1^l - C_3^l)(\overline{s_L}\gamma^\mu b_L)(\overline{\nu_\tau}\gamma_\mu\nu_l)\\ &+(C_1^l + C_3^l)(\overline{s_L}\gamma^\mu b_L)(\overline{\tau_L}\gamma_\mu l_L)\\ &+2C_3^l V_{is}(\overline{u_{Li}}\gamma^\mu b_L)(\overline{\tau_L}\gamma_\mu\nu_l)\\ &+2C_3^l V_{jb}(\overline{u_{Lj}}\gamma^\mu s_L)(\overline{\tau_L}\gamma_\mu\nu_l) + \text{h.c.}\,. \end{aligned} \tag{2.2}$$

Thus, the SMEFT Lagrangian terms that contribute to $b \to c\tau\nu$ are

$$\mathcal{L} = \left(\frac{4G_F V_{cb}\delta_{l\tau}}{\sqrt{2}} + \frac{2C_3^l V_{cs}}{\Lambda^2}\right)(\overline{c_L}\gamma^\mu b_L)(\overline{\tau_L}\gamma_\mu\nu_l). \tag{2.3}$$

We obtain:

$$\frac{R(D^{(*)})}{R(D^{(*)})^{\rm SM}} = 1 + \frac{\sqrt{2}}{G_F}\mathcal{R}e\left(\frac{V_{cs}}{V_{cb}}\frac{C_3^\tau}{\Lambda^2}\right) + \frac{\sum_{\ell=e,\mu}|C_3^\ell|^2}{2G_F^2\Lambda^4}\left|\frac{V_{cs}}{V_{cb}}\right|^2, \tag{2.4}$$

where we assume that the contribution of the term quadratic in C_3^τ is negligible compared to the term linear in C_3^τ.

Thus, to account for the $R(D^{(*)})$ puzzle by purely $b \to c\tau\nu_\ell$, $\ell = e, \mu$, we need

$$\left(\frac{\sum_{\ell=e,\mu}|C_3^\ell|^2}{\Lambda^4}\right)^{1/2} = (0.24 \pm 0.04)\,\text{TeV}^{-2} = \frac{1}{[(2.0 \pm 0.2)\,\text{TeV}]^2}. \tag{2.5}$$

On the other hand, to account for the $R(D^{(*)})$ puzzle by purely $b \to c\tau\nu_\tau$, we need

$$\frac{C_3^\tau}{\Lambda^2} = (0.046 \pm 0.016)\,\text{TeV}^{-2} \approx \frac{1}{[(4.7 \pm 0.8)\,\text{TeV}]^2}. \tag{2.6}$$

3 Bounds on C_3^ℓ

If the $R(D^{(*)})$ puzzle is accounted for by purely $b \to c\tau\bar{\nu}_\ell$, eq. (2.5) implies that we need $|C_3^\ell|/\Lambda^2 \sim 1/(2\,\text{TeV})^2$. Eq. (2.2) implies that the C_3^ℓ term contributes, via four fermi operators with the flavor structures $\bar{s}b\bar{\tau}\ell$ and $b\bar{s}\bar{\nu}_\tau\nu_\ell$, to various flavor changing neutral current and lepton flavor violating processes which are forbidden in the SM. In this section, we obtain the constraints from the experimental upper bounds on such processes.

JHEP07(2022)077

- $B_s \to \tau^\pm \mu^\mp$.
  The $B_s \to \tau^\pm \mu^\mp$ decay rate is given by

$$\Gamma(B_s \to \tau^+ \mu^-) = \frac{|C_1^\mu + C_3^\mu|^2}{\Lambda^4} \frac{f_{B_s}^2 m_\tau^2 m_{B_s}}{64\pi} \left(1 - \frac{m_\tau^2}{m_{B_s}^2}\right)^2 . \tag{3.1}$$

  The experimental upper bound [24],

$$\mathcal{B}(B_s \to \tau^\pm \mu^\mp) < 4.2 \times 10^{-5}, \tag{3.2}$$

  implies

$$\frac{|C_1^\mu + C_3^\mu|}{\Lambda^2} < 0.073\,\text{TeV}^{-2}. \tag{3.3}$$

- $B^+ \to K^+ \tau^+ \mu^-$.
  The $B^+ \to K^+ \tau^+ \mu^-$ branching ratio is given by

$$\mathcal{B}(B^+ \to K^+ \mu^- \tau^+) = 8.2 \times 10^{-3}\,\text{TeV}^4 \times \frac{|C_1^\mu + C_3^\mu|^2}{\Lambda^4}. \tag{3.4}$$

  The experimental upper bound [25],

$$\mathcal{B}(B^+ \to K^+ \mu^- \tau^+) < 2.8 \times 10^{-5}, \tag{3.5}$$

  implies

$$\frac{|C_1^\mu + C_3^\mu|}{\Lambda^2} < 0.058\,\text{TeV}^{-2}. \tag{3.6}$$

- $B^+ \to K^+ \tau^+ e^-$.
  The $B^+ \to K^+ \tau^+ e^-$ branching ratio is given by

$$\mathcal{B}(B^+ \to K^+ e^- \tau^+) = 8.2 \times 10^{-3}\,\text{TeV}^4 \times \frac{|C_1^e + C_3^e|^2}{\Lambda^4}. \tag{3.7}$$

  The experimental upper bound [25],

$$\mathcal{B}(B^+ \to K^+ e^- \tau^+) < 1.5 \times 10^{-5}, \tag{3.8}$$

  implies (see also ref. [26])

$$\frac{|C_1^e + C_3^e|}{\Lambda^2} < 0.044\,\text{TeV}^{-2}. \tag{3.9}$$

- $B^+ \to K^+ \bar{\nu}_\tau \nu_\ell$.
  The $B^+ \to K^+ \nu\bar{\nu}$ branching ratio, normalized to the SM rate, is given by

$$\mathcal{R}_{K\nu\bar{\nu}} \equiv \frac{\mathcal{B}(B^+ \to K^+ \nu\bar{\nu})}{\mathcal{B}^{\text{SM}}(B^+ \to K^+ \nu\bar{\nu})} = 1 + 3.5 \times 10^3\,\text{TeV}^4 \times \frac{|C_1^\mu - C_3^\mu|^2 + |C_1^e - C_3^e|^2}{\Lambda^4}. \tag{3.10}$$

  The experimental upper bound [27, 28],

$$\mathcal{B}(B^+ \to K^+ \nu\bar{\nu}) < 1.6 \times 10^{-5}, \tag{3.11}$$

  which corresponds to $R_{K\nu\bar{\nu}} \lesssim 4$, implies (see also ref. [29])

$$\frac{|C_1^\ell - C_3^\ell|}{\Lambda^2} < 0.031\,\text{TeV}^{-2}. \tag{3.12}$$

JHEP07(2022)077

From the upper bounds on $|C_1^\ell \pm C_3^\ell|$ we can obtain upper bounds on C_3^ℓ alone. Assuming that C_1^ℓ and C_3^ℓ are real, we obtain:

$$\frac{|C_3^\mu|}{\Lambda^2} < 0.044\,\text{TeV}^{-2}, \quad \frac{|C_3^e|}{\Lambda^2} < 0.037\,\text{TeV}^{-2}. \tag{3.13}$$

Comparing to the requirement of eq. (2.5), we conclude that the contributions from $b \to c\tau\nu_\ell$ ($\ell = \mu, e$) can account for, at most, 3.4% of the deviation of the central value of $R(D^{(*)})$ from the SM value.

3.1 Other processes

We here list a few processes to which the operators of eq. (2.2) contribute. Due to our specific choice of flavor structure in eq. (2.1), these contributions are suppressed and provide weak limits only. With a different flavor structure, however, they might provide relevant bounds.

- Forbidden top decays, such as $t \to q\tau\ell$ with $q = c, u$ and $\ell = \mu, e$. There are currently no upper bounds on these decay rates [30].

- Forbidden tau decays, for two of which there are experimental upper bounds [31, 32]:

$$\begin{aligned} \mathcal{B}(\tau \to \mu\pi^0) &< 1.1 \times 10^{-7}, \\ \mathcal{B}(\tau \to e\pi^0) &< 8.0 \times 10^{-8}. \end{aligned} \tag{3.14}$$

  Compared to the allowed $\tau \to \nu_\tau \bar{u} q$ ($q = d, s$), with $\mathcal{B}(\tau \to \nu\bar{u}q) \sim 0.65$, there is a CKM suppression of $|V_{ub}V_{us}|^2 \sim 7 \times 10^{-7}$. Thus, branching ratios of $\mathcal{O}(10^{-7})$ do not provide significant bounds.

- Allowed τ decays: the ratio $\mathcal{B}(\tau \to K\nu)/\mathcal{B}(\tau \to \pi\nu)$ is modified because the contribution is only to the former. Both branching ratios are measured, with [30]

$$R_{K/\pi} \equiv \Gamma(K^-\nu)/\Gamma(\pi^-\nu) = (6.44 \pm 0.09) \times 10^{-2}. \tag{3.15}$$

  The SM gives [33]

$$R_{K/\pi} = \frac{f_K^2 |V_{us}|^2}{f_\pi^2 |V_{ud}|^2} \frac{(1 - m_K^2/m_\tau^2)^2}{(1 - m_\pi^2/m_\tau^2)^2} (1 + \delta_{\text{LD}}), \tag{3.16}$$

  where $\delta_{\text{LD}} = (0.03 \pm 0.44)\%$ is the long distance correction, and $f_K/f_\pi = 1.189 \pm 0.007$. The new contribution to $\tau \to K\nu$ is CKM suppressed by $|V_{ub}/V_{us}|^2 \sim 3 \times 10^{-4}$ compared to the SM contribution, and so it is well below the uncertainty from δ_{LD} and cannot be constrained.

- Allowed D_s^+ decays: the ratio $\mathcal{B}(D_s \to \tau\nu)/\mathcal{B}(D_s \to \mu\nu)$ is modified because the contribution is only to the former. Both branching ratios are measured, with [34]

$$R_{\tau/\mu} = \frac{\mathcal{B}(D_s \to \tau\nu)}{\mathcal{B}(D_s \to \mu\nu)} = 10.73 \pm 0.88. \tag{3.17}$$

JHEP07(2022)077

The SM gives

$$R^{\rm SM}_{\tau/\mu} = \frac{m_\tau^2(1-m_\tau^2/m_{D_s}^2)^2}{m_\mu^2(1-m_\mu^2/m_{D_s}^2)^2} = 9.76. \tag{3.18}$$

Within our framework,

$$\frac{R_{\tau/\mu}}{R^{\rm SM}_{\tau/\mu}} = 1 + \frac{\sum_{\ell=e,\mu}|C_3^\ell|^2}{8G_F^2\Lambda^4}\left|\frac{V_{cb}}{V_{cs}}\right|^2. \tag{3.19}$$

Given the suppression by $(1/8)|V_{cb}/V_{cs}|^2 \sim 2\times10^{-4}$, we learn that $R_{\tau/\mu}$ does not provide a significant bound.

- Forbidden $D^0 \to \tau^\pm e^\mp$ decays. There is currently no upper bound on these rates [30]. The new contributions are strongly suppressed by several factors, compared to the leading semileptonic decay: CKM suppression by either $|V_{cb}V_{us}/V_{cs}|^2 \sim 10^{-4}$ or $|V_{ub}|^2 \sim 10^{-5}$, phase space suppression by $(1-m_\tau^2/m_{D^0}^2)^2 = 0.0084$ and annihilation suppression by $f_D^2/m_D^2 \sim 0.01$.

- Forbidden J/ψ decays [35, 36]:

$$\begin{aligned}\mathcal{B}(J/\psi \to \mu^\pm\tau^\mp) &< 2.0\times10^{-6},\\ \mathcal{B}(J/\psi \to e^\pm\tau^\mp) &< 7.5\times10^{-8}.\end{aligned} \tag{3.20}$$

Compared to the allowed $J/\psi \to \ell^+\ell^-$, with $\mathcal{B}(J/\psi\to\ell^+\ell^-)\sim 0.06$, there is a CKM suppression $|V_{cb}V_{cs}|^2 \sim 1.6\times10^{-3}$. Furthermore [17, 18],

$$\frac{\mathcal{B}(J/\psi\to\tau^\pm\ell^\mp)}{\mathcal{B}(J/\psi\to\ell^+\ell^-)} \propto \left(\frac{m_{J/\psi}}{\Lambda}\right)^4. \tag{3.21}$$

Thus, the bound on $\Lambda/\sqrt{|C_3^e|}$ is of $\mathcal{O}(6m_{J/\psi})$ and on $\Lambda/\sqrt{|C_3^\mu|}$ even weaker.

4 Collider searches

The effective operators of eq. (2.1) will contribute to the scattering process $pp \to \tau^\pm\mu^\mp X_h$, where X_h stands for final hadrons. In this section, we estimate the upper bound on $|C_1^\mu + C_3^\mu|/\Lambda^2$ that can be obtained at the LHC at present and in the future. (For related work, see [37–41] and, in particular, [42–44].)

We base our estimate on the ATLAS search for new physics in $pp \to \mu^+\mu^-$ (with up to one b-jet) at $\sqrt{s} = 13\,\text{TeV}$ with $139\,\text{fb}^{-1}$ of data [45]. (The phenomenological framework for the ATLAS analysis was suggested in ref. [46].) Ref. [45] obtains a bound $\Lambda_{\mu\mu} > 2.4\,\text{TeV}$ on the scale that suppresses dimension-six $bs\mu\mu$ contact interaction. This limit is obtained in the analysis with a b-veto ($pp \to \mu^+\mu^- + 0b$), by searching for events with high dimuon mass, $m_{\mu\mu} > 1800\,\text{GeV}$. A somewhat weaker bound, $\Lambda_{\mu\mu} > 2.0\,\text{TeV}$, is obtained in the b-tag category ($pp \to \mu^+\mu^- + 1b$) with $m_{\mu\mu} > 1600\,\text{GeV}$. We deduce from these bounds the reach of ATLAS for $\Lambda_{\tau\mu}$, the scale that suppresses dimension-six $bs\tau\mu$ contact interaction.

The bound on $\Lambda_{\mu\mu}$ is inferred from the upper bound on $\sigma_{\mu\mu}$, the $\mu\mu$ signal cross section. When comparing to it a search for $pp \to \tau_h^\pm\mu^\mp + 0b$, one muon is replaced by a hadronically

JHEP07(2022)077

decaying tau-lepton. Given that $\mathcal{B}(\tau \to \text{hadrons}) \approx 2/3$, then at similar energy and for $\Lambda_{\tau\mu} = \Lambda_{\mu\mu}$, the signal cross-sections fulfill

$$\sigma_{\tau_h\mu}/\sigma_{\mu\mu} \approx 2/3. \tag{4.1}$$

The leading SM background at the high $m_{\mu\mu}$ is the Drell-Yan process to two muons, which is suppressed in a $\tau\mu$ final state selection. As concerns the Drell-Yan process to two tau-leptons, its contribution to the background is suppressed by demanding that one of the two tau-leptons decays hadronically (with branching ratio $\sim 2/3$) and the other muonically (with branching ratio $\sim 1/6$):

$$\sigma^{\tau_h\mu}_{Z/\gamma^*}/\sigma^{\mu\mu}_{Z/\gamma^*} \approx 2 \times 2/3 \times 1/6 = 2/9. \tag{4.2}$$

The next most significant SM background is the top-quark contribution, composed of $t\bar{t}$, Wt and $Wt\bar{t}$. In each of these, the two final leptons arise from two independent decay chains, e.g. $t \to l^+$ and $\bar{t} \to l^-$ in $t\bar{t}$ events. Thus,

$$\sigma^{\tau_h\mu}_{t}/\sigma^{\mu\mu}_{t} \approx 2 \times 2/3 = 4/3. \tag{4.3}$$

Defining $r_{Z/t} \equiv \sigma^{\mu\mu}_{Z/\gamma^*}/\sigma^{\mu\mu}_{t}$, we have

$$\frac{\sigma^{\text{bkgd}}_{\tau_h\mu}}{\sigma^{\text{bkgd}}_{\mu\mu}} \approx \frac{(2/9)r_{Z/t} + 4/3}{r_{Z/t} + 1}. \tag{4.4}$$

We thus estimate

$$\frac{(s/\sqrt{b})_{\tau_h\mu}}{(s/\sqrt{b})_{\mu\mu}} \approx \frac{2}{3}\sqrt{\frac{r_{Z/t} + 1}{(2/9)r_{Z/t} + 4/3}}. \tag{4.5}$$

For $r_{Z/t} > 1$, we have $(s/\sqrt{b})_{\tau_h\mu}/(s/\sqrt{b})_{\mu\mu} \gtrsim 0.75$, which is our conservative estimate. Note that if $r_{Z/t} \gtrsim 4$, the sensitivity to $\tau\mu$ is in fact stronger than to $\mu\mu$.

Given our conservative estimate, and that the signal sensitivity is fixed by experiment, the observational significance in $\tau\mu$ would be the same as in $\mu\mu$ for $s_{\tau\mu} \sim (4/3)s_{\mu\mu}$. Since $s \propto 1/\Lambda^4$, we expect that current data can put a lower bound of

$$\Lambda_{\tau\mu} > 2.2\,\text{TeV}. \tag{4.6}$$

If the HL-LHC achieves its target integrated luminosity, $L = 4000\,\text{fb}^{-1}$, and no signal is observed, the bound would be strengthened to

$$\Lambda_{\tau\mu} > 3.3\,\text{TeV}, \tag{4.7}$$

where we scaled according to $s/\sqrt{b} \propto \sqrt{L}$ since $s/\sqrt{b} \sim L\sigma_s/\sqrt{L\sigma_b}$.

The analogous bounds $\Lambda_{ee} > 2.0(1.8)\,\text{TeV}$ are also attained in ref. [45], from the analysis of $pp \to e^+e^- + 0b(1b)$ events with high dielectron mass, $m_{ee} > 1900(1500)\,\text{GeV}$. We deduce from these bounds the reach of ATLAS for $\Lambda_{\tau e}$, the scale that suppresses dimension-six $bs\tau e$ contact interaction. The analysis goes along the same lines as the $\Lambda_{\tau\mu}$ analysis above. Given the lepton flavor universality of the relevant SM interactions, there

JHEP07(2022)077

are only negligible differences in the background cross sections. Our conservative estimates imply that current data can put a lower bound of

$$\Lambda_{\tau e} > 1.9\,\text{TeV}. \tag{4.8}$$

For the HL-LHC with $L = 4000\,\text{fb}^{-1}$, and no signal observed, the bound would be strengthened to

$$\Lambda_{\tau e} > 2.8\,\text{TeV}. \tag{4.9}$$

For the $\tau\ell + 0b$ ($\ell = \mu, e$) final states, our framework predicts additional contributions to those coming from the $b\bar{s}\tau\ell$ contact interactions. In particular, there are also contributions from $u\bar{c}\tau\ell$ and $u\bar{u}\tau\ell$ contact interactions. These contributions are, however, CKM suppressed, the former by $|V_{us}V_{cb}|^2 \sim 10^{-4}$, and the latter by $|V_{us}V_{ub}|^2 \sim 10^{-6}$. Involving the valence u-quark, instead of the sea s or b quark, gains less than two orders of magnitude from the parton distribution function, so that these contributions can be safely neglected in our analysis.

The operators of eq. (2.2) lead also to mono-τ signatures at the LHC. The ATLAS [47, 48] and CMS [49] experiments have searched for $\tau\nu$ resonances, and their results have been recasted to apply to the SMEFT operators of interest to us [50–53], yielding [51]

$$\Lambda_{\tau\nu} > 1.5\,\text{TeV}. \tag{4.10}$$

At the HL-LHC, with integrated luminosity of $L = 3000\,\text{fb}^{-1}$, the sensitivity would be strengthened to 2.3 TeV [52].

5 Conclusions

While the $R(D^{(*)})$ puzzle concerns lepton universality violation, it concerns also lepton flavor violation. Our study aimed to answer several questions in this context: can the $R(D^{(*)})$ puzzle be explained by purely lepton flavor violating new physics? Does the puzzle imply that the relevant new physics require non-generic flavor structure? Can searches for lepton flavor violation at the ATLAS and CMS experiments shed light on these questions?

We reached the following conclusions:

- Given that, to account for the central value of $R(D^{(*)})$, it is required that $|C_3^\ell|/\Lambda^2 \simeq 0.24\,\text{TeV}^{-2}$, but other constraints require that $|C_3^\mu|/\Lambda^2 < 0.044\,\text{TeV}^{-2}$, the contribution of $b \to c\tau\nu_\ell$, with $\ell = e, \mu$, to $R(D^{(*)})/R(D^{(*)})^{\text{SM}} - 1$ cannot exceed about 4% of the required shift.

- Given that, to account for the central value of $R(D^{(*)})$, it is required that $|C_3^\tau|/\Lambda^2 \simeq 0.046\,\text{TeV}^{-2}$, but phenomenological constraints require that $|C_3^\mu|/\Lambda^2 < 0.044\,\text{TeV}^{-2}$, and $|C_3^e|/\Lambda^2 \simeq 0.037\,\text{TeV}^{-2}$, we learn that no special alignment with the τ-direction is needed to explain the $R(D^{(*)})$ puzzle.

- Conversely, if operators of the form

$$\frac{C_3^l}{\Lambda^2}(\overline{L_\tau}\gamma_\sigma\tau^a L_l)(\overline{Q_s}\gamma^\sigma\tau^a Q_b) \tag{5.1}$$

JHEP07(2022)077

have C_3^τ, C_3^μ and C_3^e all of the same order of magnitude, $C_3^l/\Lambda^2 \sim 0.04\,\text{TeV}^{-2}$, then the shift in $R(D^{(*)})$ will be dominated by a factor of order 30 by C_3^τ, and all phenomenological constraints satisfied.

- Comparing eqs. (4.6) and (4.7) to eq. (3.6), we conclude that future searches at the (HL-)LHC will have to achieve an improvement in sensitivity by a factor $\sim (2.5)13$ in order to compete with existing constraints from the B-factories on $|C_1^\mu + C_3^\mu|/\Lambda^2$. Comparing eqs. (4.8) and (4.9) to eq. (3.9), we conclude that future searches at the (HL-)LHC will have to achieve an improvement in sensitivity by a factor $\sim (8.5)40$ in order to compete with existing constraints from the B-factories on $|C_1^e + C_3^e|/\Lambda^2$.

A Additional operators

In the context of $R(D^{(*)})$, the following three operators are considered, in addition to the operators of eq. (2.1):

$$\begin{aligned} O_{SR}^\ell &= (\overline{Q_s} b_R)(\overline{\tau_R} L_\ell), \\ O_{SL}^\ell &= (\overline{c_R} Q_b)(\overline{\tau_R} L_\ell), \\ O_T^\ell &= (\overline{c_R}\sigma^{\mu\nu} Q_b)(\overline{\tau_R}\sigma_{\mu\nu} L_\ell). \end{aligned} \tag{A.1}$$

At low energy, we can write the effective Hamiltonian terms that are relevant to $b \to c\tau\bar{\nu}_\ell$ transitions as follows (see e.g. [54]):

$$\mathcal{H}_{\text{eff}} = 2\sqrt{2} G_F V_{cb} [\hat{C}_S^\ell (\bar{c}b)(\overline{\tau_R}\nu_\ell) + \hat{C}_P^\ell (\bar{c}\gamma_5 b)(\overline{\tau_R}\nu_\ell) + \hat{C}_T^\ell (\overline{c_R}\sigma^{\mu\nu} b_L)(\overline{\tau_R}\sigma_{\mu\nu}\nu_\ell)], \tag{A.2}$$

where the $C_{S,P}^\ell$ coefficients are related to the $O_{SL}^\ell \pm O_{SR}^\ell$ operators. In other words, writing down the Wilson coefficients in the SMEFT as C_X/Λ^2, as in eq. (2.1), the relation with the $\hat{C}_X$ coefficients of eq. (A.2) is given by

$$\frac{C_X^\ell}{\Lambda^2} = 1.32\ \hat{C}_X^\ell\ \text{TeV}^{-2} = \frac{\hat{C}_X^\ell}{(0.87\,\text{TeV})^2}. \tag{A.3}$$

Unlike lepton flavor diagonal operators, the lepton flavor violating ones do not interfere with the SM contributions, independent of their Lorentz structure. Their contributions to various observables related to $b \to c\tau\bar{\nu}_\ell$ transitions are given by (see e.g. [54])

$$\begin{aligned} R(D)/R(D)^{\text{SM}} &= 1 + 1.09|\hat{C}_S^\ell|^2 + 0.75|\hat{C}_T^\ell|^2, \\ R(D^*)/R(D^*)^{\text{SM}} &= 1 + 0.05|\hat{C}_P^\ell|^2 + 16.3|\hat{C}_T^\ell|^2, \\ \text{BR}(B_c \to \tau\nu)/\text{BR}(B_c \to \tau\nu)^{\text{SM}} &= 1 + 18.5|\hat{C}_P^\ell|^2, \end{aligned} \tag{A.4}$$

where the Wilson coefficients C_X^ℓ are given at the scale m_b.

Accounting for the $R(D^{(*)})$ puzzle by either of these operators is disfavored compared to the one we studied above:

JHEP07(2022)077

- While $|\hat{C}_S^\ell|^2 \approx 0.13$ can account for $R(D)$, it leaves $R(D^*) = R(D^*)^{\rm SM}$.
- While $|\hat{C}_P^\ell|^2 \approx 3$ can account for $R(D^*)$, it gives $\mathrm{BR}(B_c \to \tau\nu) \approx 1$.
- While $|\hat{C}_T^\ell|^2 \approx 0.009$ can account for $R(D^*)$, it leaves $R(D) \simeq R(D)^{\rm SM}$.

A specific combination of operators may, however, account for $R(D)$ and $R(D^*)$ without violating the bound from $B_c \to \tau\nu$. The requirements are

$$|\hat{C}_S^\ell(m_b)| \approx 0.35, \quad |\hat{C}_P^\ell(m_b)| \lesssim 1.5, \quad |\hat{C}_T^\ell(m_b)| \approx 0.09. \tag{A.5}$$

This combination of parameters seems rather ad-hoc. A model that goes, however, in this direction is the R_2 leptoquark model of refs. [55, 56] ($R_2(3,2)_{+7/6}$ is a scalar, color-triplet, SU(2)-doublet of hypercharge $+7/6$). It generates at the high scale Λ of integrating out R_2 the following Wilson coefficients:

$$\hat{C}_S^\ell(\Lambda) = \hat{C}_P^\ell(\Lambda) = 4\hat{C}_T^\ell(\Lambda). \tag{A.6}$$

With $\Lambda \approx 1\,\mathrm{TeV}$, the RGE modifies this relation into (see e.g. [54])

$$\hat{C}_S^\ell(m_b) = \hat{C}_P^\ell(m_b) \approx 8.1 C_T^\ell(m_b), \tag{A.7}$$

thus predicting

$$r_{D/D^*} \equiv \frac{[R(D)/R(D)^{\rm SM}] - 1}{[R(D^*)/R(D^*)^{\rm SM}] - 1} \approx 3.7. \tag{A.8}$$

Two comments are in order:

- The original models, when fitting the data with a final ν_τ, find that $\hat{C}_S^\ell$ needs to be close to imaginary, to reduce the effect of interference terms. With a final ν_ℓ, the interference terms vanish identically, with no need for a special phase structure.
- In the absence of interference terms, the predicted ratio (A.8) is disfavored at the 3σ level by the current experimental range (we take into account the correlation between the measurements [1]):

$$r_{D/D^*} = 0.96 \pm 0.92. \tag{A.9}$$

Finally, let us mention that among the three operators of eq. (A.1), only O_{SR}^ℓ generates FCNC processes in the down sector. In particular, the contribution of $|C_P^\mu|^2$ to the $\Gamma(B_s \to \tau^+\mu^-)$ decay is enhanced by $[m_{B_s}^2/(m_\tau(m_b + m_s))]^2$ compared to the one of $|C_1^\mu + C_3^\mu|^2$. Thus the bound on $|C_P^\ell|/\Lambda^2$ is a factor $\sim m_{B_s}/m_\tau = 3$ stronger than the bound of eq. (3.3): $|C_P^\mu|/\Lambda^2 < 0.024\,\mathrm{TeV}^{-2}$ or, equivalently, $|\hat{C}_P^\mu| < 0.018$, making its contribution to $R(D^*)$ negligible.

JHEP07(2022)077

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JHEP07(2022)077