

RECEIVED: November 9, 2016

REVISED: January 23, 2017

ACCEPTED: February 5, 2017

PUBLISHED: February 13, 2017

$\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ decay in scalar and vector leptoquark scenarios

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ABSTRACT: It has been shown that the anomalies observed in $\bar{B} \rightarrow D^{(*)} \tau \bar{\nu}_\tau$ and $\bar{B} \rightarrow \bar{K} \ell^+ \ell^-$ decays can be resolved by adding a single scalar or vector leptoquark to the Standard Model, while constraints from other precision measurements in the flavour sector can be satisfied without fine-tuning. To further explore these two interesting scenarios, in this paper, we study their effects in the semi-leptonic $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ decay. Using the best-fit solutions for the operator coefficients allowed by the current data of mesonic decays, we find that (i) the two scenarios give similar amounts of enhancements to the branching fraction $\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)$ and the ratio $R_{\Lambda_c} = \mathcal{B}(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau) / \mathcal{B}(\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell)$, (ii) the two best-fit solutions in each of these two scenarios are also indistinguishable from each other, (iii) both scenarios give nearly the same predictions as those of the Standard Model for the longitudinal polarizations of Λ_c and τ as well as the lepton-side forward-backward asymmetry. With future measurements of these observables in $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ decay at the LHCb, the two leptoquark scenarios could be further tested, and even differentiated from the other NP explanations for the $R_{D^{(*)}}$ anomalies. We also discuss the feasibility for the measurements of these observables at the LHC and the future e^+e^- colliders.

KEYWORDS: Beyond Standard Model, Heavy Quark Physics

ARXIV EPRINT: [1611.01635](https://arxiv.org/abs/1611.01635)

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1 Introduction

While no direct evidences for physics beyond the Standard Model (SM) have been found at the LHC so far, there are some interesting indirect hints for New Physics (NP) in the flavour sector [1–3]. It is particularly interesting to note that intriguing effects of lepton favour universality violation (LFUV) have been observed in rare B -meson decays. To be more specific, the ratios of charged-current decays, $R_D = \frac{\mathcal{B}(\bar{B} \rightarrow D \tau \bar{\nu}_\tau)}{\mathcal{B}(\bar{B} \rightarrow D \ell \bar{\nu}_\ell)}$ and $R_{D^*} = \frac{\mathcal{B}(\bar{B} \rightarrow D^* \tau \bar{\nu}_\tau)}{\mathcal{B}(\bar{B} \rightarrow D^* \ell \bar{\nu}_\ell)}$, with $\ell = e, \mu$, have been measured by the BaBar [4, 5], Belle [6–8] and LHCb [9] Collaborations. The latest averages by the Heavy Flavor Average Groups (HFAG) [10],

$$R_D^{\text{exp.}} = 0.397 \pm 0.040(\text{stat.}) \pm 0.028(\text{syst.}), \quad R_{D^*}^{\text{exp.}} = 0.316 \pm 0.016(\text{stat.}) \pm 0.010(\text{syst.}), \quad (1.1)$$

exceed the SM predictions,

$$R_D^{\text{SM}} = 0.300 \pm 0.008 [11], \quad R_{D^*}^{\text{SM}} = 0.252 \pm 0.003 [12], \quad (1.2)$$

by 1.9σ and 3.3σ , respectively. Once the measurement correlations between R_D and R_{D^*} are taken into account, the deviation will be at 4.0σ level. Another hint of LFUV has also been reported in the $b \rightarrow s \ell^+ \ell^-$ process by the LHCb experiment [13]:

$$R_K^{\text{exp.}} = \frac{\Gamma(B^+ \rightarrow K^+ \mu^+ \mu^-)}{\Gamma(B^+ \rightarrow K^+ e^+ e^-)} \Bigg|_{q^2 \in [1,6] \text{ GeV}^2} = 0.745_{-0.074}^{+0.090}(\text{stat.}) \pm 0.036(\text{syst.}), \quad (1.3)$$

which is about 2.6σ lower than the corresponding SM prediction $R_K^{\text{SM}} = 1.00 \pm 0.03$ [14, 15].

The observed $R_{D^{(*)}}$ and R_K anomalies, if confirmed with future more precise data, would be clear signs for NP beyond the SM, and have already inspired lots of studies; for a recent review, the readers are referred to refs. [1–3] and references therein. Here we are interested in the possible NP solutions with a single scalar or vector leptoquark (LQ) scenario [16, 17]. In ref. [16], it has been shown that the anomalies $R_{D^{(*)}}$, R_K and $(g-2)_\mu$ could be addressed by adding to the SM just one TeV-scale scalar LQ transforming as $(\mathbf{3}, \mathbf{1}, -\frac{1}{3})$ under the SM gauge group. On the other hand, as shown in ref. [17], the $R_{D^{(*)}}$, R_K and the angular observable P'_5 in $\bar{B} \rightarrow \bar{K}^* \mu^+ \mu^-$ decay could be explained by just one vector LQ transforming as $(\mathbf{3}, \mathbf{3}, \frac{2}{3})$ under the SM gauge group. Under the constraints from both the ratios $R_{D^{(*)}}$ and the q^2 spectra of $\bar{B} \rightarrow D^{(*)} \tau \bar{\nu}_\tau$ decays provided by the BaBar [5] and Belle [6, 7] Collaborations, four best-fit solutions are found for the operator coefficients induced by the scalar LQ [18], two of which are, however, already excluded by the purely leptonic $B_c^- \rightarrow \tau^- \bar{\nu}_\tau$ decay [19]. At the same time, two best-fit solutions are also found for the operator coefficients induced by the vector LQ [18].

To further explore the two interesting LQ scenarios, in this paper, we shall study their effects in the semi-leptonic $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ decay, which is induced by the same quark-level transition as the $\bar{B} \rightarrow D^{(*)} \tau \bar{\nu}_\tau$ decays. While the Λ_b baryons are not produced at an e^+e^- B -factory, they account for around 20% of the b -hadrons produced at the LHC [20]. Remarkably, the produced number of Λ_b baryons is comparable to that of B_u or B_d mesons, and is significantly higher than that of B_s meson [20, 21]. Due to the spin-half nature of Λ_b , its decay may provide complementary information compared to the corresponding mesonic one. Motivated by the $R_{D^{(*)}}$ anomalies, the semi-leptonic $\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell$ decay has been studied recently in refs. [22–27].

In this paper, besides the total and differential branching fractions, as well as the ratio $R_{\Lambda_c} = \frac{\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)}{\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell)}$, discussed already in previous studies [23–26], we shall also discuss the longitudinal polarizations of the daughter baryon Λ_c and the τ lepton, and the lepton-side forward-backward asymmetry in this decay [24]. The feasibility for the measurements of these observables at the LHC, which is the currently available experiment to explore the Λ_b decays, as well as at the future e^+e^- colliders, such as the International Linear Collider (ILC) and the Circular Electron Positron Collider (CEPC), will also be discussed. We calculate these observables using the helicity formalism developed in refs. [28, 29], and have rederived and confirmed the helicity amplitudes associated with both the (axial-)vector and (pseudo-)scalar interactions given already in refs. [23–25]; for the (pseudo-)tensor-type current, however, the corresponding helicity amplitudes are new and presented here for the first time. As the $\Lambda_b \rightarrow \Lambda_c$ transition form factors are not yet determined quite well and still bring large uncertainties, it is instructive to check the sensitivity of these observables to the different values of form factors obtained, for example, in a covariant confined quark model [24] (used in ref. [24]), in the QCD sum rules [30] (used in ref. [23]), or in the lattice calculations [26] (used in refs. [25, 26]). To this end, rather than choosing a single form for these form factors, we use, as a comparison, the results obtained both from the QCD sum rules [30], which satisfy the heavy quark effective theory (HQET) relations [31–33], and from the latest lattice calculations with $2+1$ dynamical flavours [26].¹ Using the best-fit

¹There are currently only the lattice results for the (axial-)vector form factors [26]. For the (pseudo-)tensor form factors, since there are no lattice results yet, we still use the HQET relations to relate them to the corresponding (axial-)vector ones.

solutions for the operator coefficients allowed by the current data of mesonic decays, we find that the two scenarios give similar amounts of enhancements relative to the SM predictions for the branching fraction $\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)$ and the ratio R_{Λ_c} , and the two best-fit solutions in each of the two scenarios are also indistinguishable from each other based only on these two observables. On the other hand, both of these two scenarios give nearly the same predictions as the SM for the longitudinal polarizations of Λ_c and τ as well as the lepton forward-backward asymmetry. With future precise measurements of these observables at the LHCb, the two scenarios could be further tested and even differentiated from the other explanations to the $R_{D^{(*)}}$ anomalies [18, 34–94].

This paper is organized as follows: in section 2, we recapitulate briefly both the scalar and vector LQ scenarios [16, 17]. In section 3, we calculate the helicity amplitudes and list the relevant physical observables for the semi-leptonic $\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell$ decays. In section 4, the scalar and vector LQ effects on the branching fraction $\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)$, the ratio R_{Λ_c} , the Λ_c and τ longitudinal polarizations, as well as the lepton-side forward-backward asymmetry are discussed. We finally conclude in section 5. The $\Lambda_b \rightarrow \Lambda_c$ transition form factors and the helicity-dependent differential decay rates are collected in appendices A and B, respectively.

2 The scalar and vector LQ scenarios

In this section, we recapitulate the scalar and vector LQ models, where a single TeV-scale scalar or vector LQ is added to the SM to address the aforementioned anomalies [16, 17]. For a recent comprehensive review of LQ models, the readers are referred to ref. [95].

2.1 The scalar LQ scenario

Firstly, we consider the scalar LQ ϕ transforming as $(\mathbf{3}, \mathbf{1}, -\frac{1}{3})$ under the SM gauge group, in which its couplings to SM fermions are described by the Lagrangian [16]

$$\mathcal{L}_{\text{int}}^\phi = \bar{Q}_L^c \boldsymbol{\lambda}^L i\tau_2 L\phi^* + \bar{u}_R^c \boldsymbol{\lambda}^R \ell_R \phi^* + \text{h.c.}, \quad (2.1)$$

where $\boldsymbol{\lambda}^{L,R}$ are the Yukawa coupling matrices in flavour space, and Q_L, L denote the left-handed quark and lepton doublet, while u_R, ℓ_R the right-handed up-type quark and lepton singlet, respectively. The charge-conjugated spinors are defined as $\psi^c = C\bar{\psi}^T$, $\bar{\psi}^c = \psi^T C$ ($C = i\gamma^2\gamma^0$). Such a scalar ϕ mediates the $b \rightarrow c\tau\bar{\nu}_\tau$ decay at tree level, and the resulting effective weak Hamiltonian including the SM contribution is given as [16, 19]

$$\mathcal{H}_{\text{eff}} = \frac{4G_F V_{cb}}{\sqrt{2}} \left[C_V \bar{c}\gamma_\mu P_L b \bar{\tau}\gamma^\mu P_L \nu_\tau + C_S \bar{c}P_L b \bar{\tau}P_L \nu_\tau - \frac{1}{4}C_T \bar{c}\sigma_{\mu\nu} P_L b \bar{\tau}\sigma^{\mu\nu} P_L \nu_\tau \right], \quad (2.2)$$

where C_V, C_S, C_T are the Wilson coefficients of the corresponding four-fermion operators and, at the matching scale $\mu = M_\phi$, are given explicitly as

$$C_V(M_\phi) = 1 + \frac{\lambda_{b\nu_\tau}^L \lambda_{c\tau}^{L*}}{4\sqrt{2}G_F V_{cb} M_\phi^2}, \quad (2.3)$$

$$C_S(M_\phi) = C_T(M_\phi) = -\frac{\lambda_{b\nu_\tau}^L \lambda_{c\tau}^{R*}}{4\sqrt{2}G_F V_{cb} M_\phi^2}. \quad (2.4)$$

In order to resum potentially large logarithmic effects, the Wilson coefficients C_S and C_T should be run down to the characteristic scale of the process we are interested in, i.e., $\mu_b \sim m_b$, while C_V is not renormalized because of the conservation of vector currents. The explicit evolution equations could be found, for example, in ref. [95].

As shown in ref. [16], such a scalar LQ could explain the $R_{D^{(*)}}$, R_K and $(g-2)_\mu$ anomalies, while constraints from other precision measurements in the flavour sector can be satisfied without fine-tuning. Especially, under the constraints from both the ratios $R_{D^{(*)}}$ and the measured q^2 spectra in $B \rightarrow D^{(*)} \tau \bar{\nu}_\tau$ decays, four best-fit solutions are found for the operator coefficients induced by the scalar LQ [18], two of which are, however, already excluded by the purely leptonic $B_c^- \rightarrow \tau^- \bar{\nu}_\tau$ decay [19]. Consequently, in this paper, we shall consider only the remaining two solutions denoted by P_A and P_C in ref. [19].

2.2 The vector LQ scenario

We now introduce the second scenario in which the SM is extended by a vector $SU(2)_L$ triplet U_3^μ transforming as $(\mathbf{3}, \mathbf{3}, \frac{2}{3})$ under the SM gauge group. The coupling of the vector multiplet U_3^μ to a lepton-to-quark current with $(V-A)$ structure is given by [17]

$$\mathcal{L}_{U_3} = g_{ij} \bar{Q}_i \gamma^\mu \tau^A U_{3\mu}^A L_j + \text{h.c.}, \quad (2.5)$$

where τ^A ($A = 1, 2, 3$) are the Pauli matrices in the $SU(2)_L$ space, and L_i and Q_i ($i, j = 1, 2, 3$) denote the left-handed lepton and quark doublets, respectively. The Lagrangian eq. (2.5) is written in the fermion mass basis, with g_{ij} defined as the couplings of the $Q = 2/3$ component of the triplet, $U_{3\mu}^{(2/3)}$, to \bar{d}_{Li} and ℓ_{Lj} . Expanding the $SU(2)_L$ components, we get explicitly

$$\begin{aligned} \mathcal{L}_{U_3} = & U_{3\mu}^{(2/3)} [\mathcal{V}g\mathcal{U}]_{ij} \bar{u}_i \gamma^\mu P_L \nu_j - g_{ij} \bar{d}_i \gamma^\mu P_L \ell_j \\ & + U_{3\mu}^{(5/3)} (\sqrt{2}\mathcal{V}g)_{ij} \bar{u}_i \gamma^\mu P_L \ell_j \\ & + U_{3\mu}^{(-1/3)} (\sqrt{2}g\mathcal{U})_{ij} \bar{d}_i \gamma^\mu P_L \nu_j + \text{h.c.}, \end{aligned} \quad (2.6)$$

where \mathcal{V} and \mathcal{U} represent the Cabibbo-Kobayashi-Maskawa (CKM) [96, 97] and the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) [98, 99] matrix, respectively. Here we assume the neutrinos to be massless and, therefore, the PMNS matrix can be rotated away through field redefinitions.

The vector multiplet U_3^μ can also mediate the $b \rightarrow c \tau \bar{\nu}_\tau$ transitions at tree level, and the resulting effective weak Hamiltonian including the SM contribution can be written as [17]

$$\mathcal{H}_{\text{eff}} = \frac{4G_F V_{cb}}{\sqrt{2}} C'_V (\bar{c} \gamma^\mu P_L b) (\bar{\tau} \gamma_\mu P_L \nu_\tau), \quad (2.7)$$

where C'_V is the Wilson coefficient at the matching scale $\mu = M_U$ and is given by

$$C'_V = 1 + \frac{\sqrt{2}g_{b\tau}^* (\mathcal{V}g)_{c\tau}}{4G_F V_{cb} M_U^2}. \quad (2.8)$$

Unlike in the scalar LQ case, the vector LQ only generates $(V-A)$ couplings and, therefore, the Wilson coefficient C'_V need not be renormalized.

As shown in ref. [17], the vector LQ scenario could also accommodate the $R_{D^{(*)}}$, R_K as well as the angular observable P_5' in $\bar{B} \rightarrow \bar{K}^* \mu^+ \mu^-$ decay. Fitting to the measured ratios $R_{D^{(*)}}$, along with acceptable q^2 spectra, two best-fit solutions, denoted as R_A and R_B , respectively, are found in this scenario [18]:

$$g_{b\tau}^*(\mathcal{V}g)_{c\tau} = \begin{cases} 0.18 \pm 0.04, R_A \\ -2.88 \pm 0.04, R_B \end{cases}, \quad (2.9)$$

where $M_U = 1 \text{ TeV}$ is taken as a benchmark. It should be noted that the triplet nature of U_3^μ also leads to various charged lepton-flavour-violating decays, such as the $B \rightarrow K \mu \tau$ and $\Upsilon(nS) \rightarrow \tau \mu$ decays, which have been discussed in refs. [17, 88].

3 $\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell$ decays in scalar and vector LQ scenarios

3.1 Helicity amplitudes

In this subsection, we give the helicity amplitudes for the process $\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell$ both within the SM and in the two LQ scenarios. Following refs. [24, 100] and starting with the effective weak Hamiltonian given by eqs. (2.2) and (2.7), one can get the helicity amplitudes of the decay. Since all types of the leptonic helicity amplitudes can be found in ref. [44], we give only the hadronic helicity amplitudes. For the $(V - A)$ -type current, we have [24]

$$H_{\lambda_2, \lambda_W} = H_{\lambda_2, \lambda_W}^V - H_{\lambda_2, \lambda_W}^A, \quad H_{\lambda_2, \lambda_W}^{V(A)} = \epsilon^{\dagger\mu}(\lambda_W) \langle \Lambda_c, \lambda_2 | \bar{c} \gamma_\mu (\gamma_\mu \gamma_5) b | \Lambda_b, \lambda_1 \rangle, \quad (3.1)$$

where λ_2 and λ_W denote the helicities of the daughter baryon Λ_c and the effective (axial-)vector-type current, respectively. The explicit expressions of H_{λ_2, λ_W} in terms of the hadronic matrix elements defined by eqs. (A.1) and (A.2) could be found in ref. [24]. For the $(S - P)$ -type current, the corresponding helicity amplitudes are given by [23]

$$H_{\lambda_2, 0}^{SP} = H_{\lambda_2, 0}^S - H_{\lambda_2, 0}^P, \quad (3.2)$$

$$H_{\pm\frac{1}{2}, 0}^{SP} = \frac{\sqrt{Q_+}}{m_b - m_c} \left(F_1^V M_- + F_3^V \frac{q^2}{M_1} \right) \pm \frac{\sqrt{Q_-}}{m_b + m_c} \left(F_1^A M_+ - F_3^A \frac{q^2}{M_1} \right), \quad (3.3)$$

where we use the abbreviations $M_\pm = M_{\Lambda_b} \pm M_{\Lambda_c}$ and $Q_\pm = M_\pm^2 - q^2$. The hadronic helicity amplitudes of the (pseudo-)tensor-type current are defined as

$$H_{\lambda_2, \lambda_W, \lambda_{W'}}^T = \epsilon^{\dagger\mu}(\lambda_W) \epsilon^{\dagger\nu}(\lambda_{W'}) \langle \Lambda_c, \lambda_2 | \bar{c} i \sigma_{\mu\nu} (1 - \gamma_5) b | \Lambda_b, \lambda_1 \rangle, \quad (3.4)$$

and their explicit expressions, in terms of the hadronic matrix elements defined by eqs. (A.5) and (A.6), are given, respectively, by

$$\begin{aligned} H_{\frac{1}{2}, +, 0}^T &= -\sqrt{\frac{2}{q^2}} \left(f_T \sqrt{Q_+} M_- + g_T \sqrt{Q_-} M_+ \right), \\ H_{\frac{1}{2}, +, -}^T &= -f_T \sqrt{Q_+} - g_T \sqrt{Q_-}, \\ H_{\frac{1}{2}, +, t}^T &= -\sqrt{\frac{2}{q^2}} \left(f_T \sqrt{Q_-} M_+ + g_T \sqrt{Q_+} M_- \right) + \sqrt{2q^2} \left(f_T^V \sqrt{Q_-} - g_T^V \sqrt{Q_+} \right), \end{aligned}$$

$$\begin{aligned}
 H_{\frac{1}{2},0,t}^T &= -f_T\sqrt{Q_-} - g_T\sqrt{Q_+} + f_T^V\sqrt{Q_-}M_+ \\
 &\quad - g_T^V\sqrt{Q_+}M_- + f_T^S\sqrt{Q_-}Q_+ + g_T^S\sqrt{Q_+}Q_-, \\
 H_{-\frac{1}{2},+,-}^T &= f_T\sqrt{Q_+} - g_T\sqrt{Q_-}, \\
 H_{-\frac{1}{2},0,-}^T &= \sqrt{\frac{2}{q^2}} \left(f_T\sqrt{Q_+}M_- - g_T\sqrt{Q_-}M_+ \right), \\
 H_{-\frac{1}{2},0,t}^T &= -f_T\sqrt{Q_-} + g_T\sqrt{Q_+} + f_T^V\sqrt{Q_-}M_+ \\
 &\quad + g_T^V\sqrt{Q_+}M_- + f_T^S\sqrt{Q_-}Q_+ - g_T^S\sqrt{Q_+}Q_-, \\
 H_{-\frac{1}{2},-,t}^T &= -\sqrt{\frac{2}{q^2}} \left(f_T\sqrt{Q_-}M_+ - g_T\sqrt{Q_+}M_- \right) + \sqrt{2q^2} \left(f_T^V\sqrt{Q_-} + g_T^V\sqrt{Q_+} \right). \quad (3.5)
 \end{aligned}$$

The helicity amplitudes satisfy the relations $H_{\lambda_2, \lambda_W, \lambda_W}^T = 0$ and $H_{\lambda_2, \lambda_W, \lambda_{W'}}^T = -H_{\lambda_2, \lambda_{W'}, \lambda_W}^T$, while all the others are found to be zero. Using the HQET relations given by eq. (A.7), we can further simplify these helicity amplitudes.

3.2 Observables in $\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell$ decays

Here we follow the conventions used in refs. [23–25], and write the two-fold differential angular decay distribution as

$$\frac{d^2\Gamma(\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell)}{dq^2 d\cos\theta_\ell} = N \left[A_1 + \frac{m_\ell^2}{q^2} (A_2^V + A_2^T) + 2A_3 + \frac{4m_\ell}{\sqrt{q^2}} A_4 + A_5 \right], \quad (3.6)$$

with

$$\begin{aligned}
 N &= \frac{G_F^2 |V_{cb}|^2 q^2 |\vec{\mathbf{p}}_2|}{512\pi^3 M_{\Lambda_b}^2} \left(1 - \frac{m_\ell^2}{q^2} \right)^2, \\
 A_1 &= C_V^2 \left[2\sin^2\theta_\ell (H_{\frac{1}{2},0}^2 + H_{-\frac{1}{2},0}^2) + (1 - \cos\theta_\ell)^2 H_{\frac{1}{2},+}^2 + (1 + \cos\theta_\ell)^2 H_{-\frac{1}{2},-}^2 \right], \\
 A_2^V &= C_V^2 \left[2\cos^2\theta_\ell (H_{\frac{1}{2},0}^2 + H_{-\frac{1}{2},0}^2) + \sin^2\theta_\ell (H_{\frac{1}{2},+}^2 + H_{-\frac{1}{2},-}^2) + 2(H_{\frac{1}{2},t}^2 + H_{-\frac{1}{2},t}^2) \right. \\
 &\quad \left. - 4\cos\theta_\ell (H_{\frac{1}{2},0} H_{\frac{1}{2},t} + H_{-\frac{1}{2},0} H_{-\frac{1}{2},t}) \right], \\
 A_2^T &= \frac{C_T^2}{4} \left[2\sin^2\theta_\ell (H_{\frac{1}{2},+,-}^{T2} + H_{\frac{1}{2},0,t}^{T2} + H_{-\frac{1}{2},+,-}^{T2} \right. \\
 &\quad \left. + H_{-\frac{1}{2},0,t}^{T2} + 2H_{\frac{1}{2},+,-}^T H_{\frac{1}{2},0,t}^T - H_{\frac{1}{2},0,t}^T H_{-\frac{1}{2},+,-}^T + 2H_{-\frac{1}{2},+,-}^T H_{-\frac{1}{2},0,t}^T \right. \\
 &\quad \left. + (1 + \cos\theta_\ell)^2 (H_{-\frac{1}{2},0,-}^{T2} + H_{-\frac{1}{2},-,t}^{T2} + 2H_{-\frac{1}{2},0,-}^T H_{-\frac{1}{2},-,t}^T) \right. \\
 &\quad \left. + (1 - \cos\theta_\ell)^2 (H_{\frac{1}{2},+,0}^{T2} + H_{\frac{1}{2},+,t}^{T2} + 2H_{\frac{1}{2},+,0}^T H_{\frac{1}{2},+,t}^T) \right], \\
 A_3 &= \frac{C_T^2}{8} \left[2\cos^2\theta_\ell (H_{\frac{1}{2},+,-}^{T2} + H_{\frac{1}{2},0,t}^{T2} + H_{-\frac{1}{2},+,-}^{T2} + H_{-\frac{1}{2},0,t}^{T2} \right. \\
 &\quad \left. + 2H_{\frac{1}{2},+,-}^T H_{\frac{1}{2},0,t}^T - H_{\frac{1}{2},0,t}^T H_{-\frac{1}{2},+,-}^T + 2H_{-\frac{1}{2},+,-}^T H_{-\frac{1}{2},0,t}^T) \right. \\
 &\quad \left. + \sin^2\theta_\ell (H_{\frac{1}{2},+,0}^{T2} + H_{\frac{1}{2},+,t}^{T2} + H_{-\frac{1}{2},0,-}^{T2} + H_{-\frac{1}{2},-,t}^{T2} \right. \\
 &\quad \left. + 2H_{\frac{1}{2},+,0}^T H_{\frac{1}{2},+,t}^T + 2H_{-\frac{1}{2},0,-}^T H_{-\frac{1}{2},-,t}^T) \right] \\
 &\quad + C_S^2 \left(H_{\frac{1}{2},0}^{SP2} + H_{-\frac{1}{2},0}^{SP2} \right),
 \end{aligned}$$

$$\begin{aligned}
 A_4 = & C_V C_S \left[-\cos\theta_\ell (H_{\frac{1}{2},0} H_{\frac{1}{2},0}^{SP} + H_{-\frac{1}{2},0} H_{-\frac{1}{2},0}^{SP}) + (H_{\frac{1}{2},t} H_{\frac{1}{2},0}^{SP} + H_{-\frac{1}{2},t} H_{-\frac{1}{2},0}^{SP}) \right] \\
 & + C_V C_T \left[\frac{\cos^2\theta}{2} (H_{\frac{1}{2},0} H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0} H_{\frac{1}{2},0,t}^T + H_{-\frac{1}{2},0} H_{-\frac{1}{2},+,-}^T + H_{-\frac{1}{2},0} H_{-\frac{1}{2},0,t}^T) \right. \\
 & - \frac{\cos\theta}{2} (H_{\frac{1}{2},t} H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},t} H_{\frac{1}{2},0,t}^T + H_{-\frac{1}{2},t} H_{-\frac{1}{2},+,-}^T + H_{-\frac{1}{2},t} H_{-\frac{1}{2},0,t}^T) \\
 & + \frac{(1-\cos\theta)^2}{4} (H_{\frac{1}{2},+} H_{\frac{1}{2},+,0}^T + H_{\frac{1}{2},+} H_{\frac{1}{2},+,t}^T) \\
 & + \frac{(1+\cos\theta)^2}{4} (H_{-\frac{1}{2},-} H_{-\frac{1}{2},0,-}^T + H_{-\frac{1}{2},-} H_{-\frac{1}{2},-,t}^T) \\
 & + \frac{\sin^2\theta}{4} (H_{\frac{1}{2},+} H_{\frac{1}{2},+,0}^T + H_{\frac{1}{2},+} H_{\frac{1}{2},+,t}^T + H_{-\frac{1}{2},-} H_{-\frac{1}{2},0,-}^T + H_{-\frac{1}{2},-} H_{-\frac{1}{2},-,t}^T \\
 & \quad \left. + 2H_{\frac{1}{2},0} H_{\frac{1}{2},+,-}^T + 2H_{\frac{1}{2},0} H_{\frac{1}{2},0,t}^T + 2H_{-\frac{1}{2},0} H_{-\frac{1}{2},+,-}^T + 2H_{-\frac{1}{2},0} H_{-\frac{1}{2},0,t}^T) \right], \\
 A_5 = & -2C_S C_T \cos\theta \left(H_{\frac{1}{2},0} H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0} H_{\frac{1}{2},0,t}^T + H_{-\frac{1}{2},0} H_{-\frac{1}{2},+,-}^T + H_{-\frac{1}{2},0} H_{-\frac{1}{2},0,t}^T \right), \tag{3.7}
 \end{aligned}$$

where $|\vec{\mathbf{p}}_2| = \sqrt{Q_+ Q_-} / (2M_{\Lambda_b})$ is the Λ_c momentum in the Λ_b rest frame, q^2 the momentum transfer squared, and θ_ℓ the polar angle of the lepton, as defined in figure 1 of ref. [24]. Integrating out $\cos\theta_\ell$ in eq. (3.6), one can then obtain the differential decay rate $d\Gamma(\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell) / dq^2$. The above results are given for the scalar LQ scenario. For the vector LQ case, we need only to replace C_V by C'_V given by eq. (2.8), while setting C_S and C_T to zero.

With eqs. (3.6) and (3.7) at hand, we can get the following physical observables:

- The differential and total branching fractions

$$\frac{d\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell)}{dq^2} = \tau_{\Lambda_b} \frac{d\Gamma(\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell)}{dq^2}, \quad \mathcal{B}(\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell) = \int_{m_\ell^2}^{M_-^2} dq^2 \frac{d\mathcal{B}}{dq^2}, \tag{3.8}$$

where τ_{Λ_b} is the lifetime of Λ_b baryon, and m_ℓ the lepton mass.

- The differential and integrated ratios

$$R_{\Lambda_c}(q^2) = \frac{d\Gamma(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau) / dq^2}{d\Gamma(\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell) / dq^2}, \quad R_{\Lambda_c} = \frac{\int_{m_\tau^2}^{M_-^2} dq^2 d\Gamma(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau) / dq^2}{\int_{m_\ell^2}^{M_-^2} dq^2 d\Gamma(\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell) / dq^2}. \tag{3.9}$$

- The lepton-side forward-backward asymmetry

$$A_{\text{FB}}(q^2) = \frac{\int_0^1 d\cos\theta (d^2\Gamma/dq^2 d\cos\theta) - \int_{-1}^0 d\cos\theta (d^2\Gamma/dq^2 d\cos\theta)}{d\Gamma/dq^2}, \tag{3.10}$$

defined as the relative difference between the differential decay rates where the angle θ_ℓ is smaller or greater than $\pi/2$.

Once the individual helicity-dependent differential decay rates are calculated, which are collected in appendix B, we can obtain another two observables, the q^2 -dependent longitudinal polarizations of Λ_c baryon and τ lepton, which are defined, respectively, as

$$P_L^{\Lambda_c}(q^2) = \frac{d\Gamma^{\lambda_2=1/2}/dq^2 - d\Gamma^{\lambda_2=-1/2}/dq^2}{d\Gamma/dq^2}, \quad P_L^\tau(q^2) = \frac{d\Gamma^{\lambda_\tau=1/2}/dq^2 - d\Gamma^{\lambda_\tau=-1/2}/dq^2}{d\Gamma/dq^2}. \quad (3.11)$$

Although the case with a polarized Λ_b can bring in a number of new observables [101, 102], we assume here that the parent baryon Λ_b is unpolarized, based on the observation that the Λ_b polarization in the LHCb setup is measured to be small and compatible with zero [103]. This leaves with us the above interesting observables. Unlike the case for light leptons $\ell = e, \mu$, in which the phase space can be fully constrained [104], the decay $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ poses several experimental challenges. Firstly, it is not possible to determine the Λ_b momentum from the tagging side at the LHCb. Secondly, as the τ decays inside the detector, there exist at least two neutrinos in the final state, prohibiting a direct signal-side reconstruction. Finally, the fact that the Λ_c is long lived and decays only weakly increases further the difficulty in determining the Λ_b decay vertex. Thus, it is very challenging to determine the kinematic distributions on an event-by-event basis at the LHCb. However, the heuristic methods employed to determine approximately the B -meson momentum in $B \rightarrow D^* \tau \bar{\nu}_\tau$ decay [9] open the possibility to study the b -hadron decays with multiple missing particles in hadron colliders. Also, it might be feasible to measure this decay at the future e^+e^- colliders, such as the ILC and CEPC, where two jets of hadrons are produced at the Z^0 peak and appears mostly in opposite side with a large boost. These facilities are also featured by the high reconstruction and tagging efficiencies, as well as the capability to measure the missing momentum.²

As detailed recently by Ivanov, Körner and Tranin in ref. [94], the information on the τ polarization can be extracted from the angular distribution of its subsequent decay modes, such as the hadronic $\tau \rightarrow \pi \nu_\tau$ and $\tau \rightarrow \rho \nu_\tau$ and the leptonic $\tau \rightarrow \mu \bar{\nu}_\mu \nu_\tau$ and $\tau \rightarrow e \bar{\nu}_e \nu_\tau$ decays. Especially, the analyzing power of the decay $\tau \rightarrow \pi \nu_\tau$ is found to be 100% [94, 105]. It is particularly interesting to note that the Belle Collaboration has recently reported on the first measurement of the τ longitudinal polarization in the decay $B \rightarrow D^* \tau \bar{\nu}_\tau$ with the subsequent decays $\tau \rightarrow \pi \nu_\tau$ and $\tau \rightarrow \rho \nu_\tau$ [106]. Although the experimental environment is different, this pioneering measurement would be very beneficial for future detailed studies of the τ polarization at the LHC and future e^+e^- colliders.

The Λ_c polarization can also be probed by analyzing the angular decay distribution of its subsequent decays, among which the $\Lambda_c \rightarrow \Lambda \pi$ and $\Lambda_c \rightarrow \Lambda \ell \nu_\ell$ modes are of particular interest [24, 107, 108]. As shown in ref. [24], the analyzing power of $\Lambda_c \rightarrow \Lambda \pi$ is close to maximal, and that of $\Lambda_c \rightarrow \Lambda \ell \nu_\ell$, although being not as large as the former, may still lead to reliable measurement. The biggest experimental challenge for the Λ_c polarization measurement in the cascade decay $\Lambda_b \rightarrow \Lambda_c (\rightarrow \Lambda \pi) \ell \bar{\nu}_\ell$ is still how to reconstruct the Λ_b

²It should be mentioned that, even if measured, the q^2 distributions will be sculpted by the phase space cuts; for example, in practice one cannot integrate over all the helicity angle θ_ℓ , because minimum lepton-energy cuts might be needed to somewhat isolate the Λ_c baryon and the lepton.

Parameter	Value	Reference
G_F	$1.166378 \times 10^{-5} \text{ GeV}^{-2}$	[113]
$\alpha_s(M_Z)$	0.1185 ± 0.0006	[113]
M_Z	91.188 GeV	[113]
m_t	$(173.21 \pm 0.87) \text{ GeV}$	[113]
$m_b(m_b)$	$(4.18 \pm 0.03) \text{ GeV}$	[113]
$m_c(m_c)$	$(1.275 \pm 0.025) \text{ GeV}$	[113]
τ_{Λ_b}	1.466 ps	[113]
M_{Λ_b}	5.61951 GeV	[113]
M_{Λ_c}	2.28646 GeV	[113]
m_τ	1.7769 GeV	[113]
m_μ	105.66 MeV	[113]
m_e	0.511 MeV	[113]
$ V_{cb} $	$(41.1 \pm 1.3) \times 10^{-3}$	[113]

Table 1. Input parameters used in our numerical analyses.

continuum model	κ	$F_1^V(q^2)$	$-F_2^V(q^2)/M_{\Lambda_b}$
rectangular	1	$6.66/(20.27 - q^2)$	$-0.21/(15.15 - q^2)$
rectangular	2	$8.13/(22.50 - q^2)$	$-0.22/(13.63 - q^2)$
triangular	3	$13.74/(26.68 - q^2)$	$-0.41/(18.65 - q^2)$
triangular	4	$16.17/(29.12 - q^2)$	$-0.45/(19.04 - q^2)$

Table 2. Pole parameterizations of the $\Lambda_b \rightarrow \Lambda_c$ transition form factors for two values of κ and for two choices of the continuum model in QCD sum rules [30].

rest frame mentioned already. But it is still hoped that, with more and more information on the Λ_c decays from the Belle [109] and BESIII [110–112] Collaborations, and sufficient statistics for the Λ_b baryon at the LHC and future e^+e^- colliders, the Λ_c polarization could be measured in this decay.

4 Numerical results and discussions

4.1 Input parameters

In this section, we investigate the scalar and vector LQ effects on the aforementioned observables, to see if their effects are large enough to cause sizable deviations from the corresponding SM predictions. Firstly, we collect in table 1 all the input parameters used in this paper.

κ	$\mathcal{B}(\Lambda_b \rightarrow \Lambda_c e \bar{\nu}_e)$	$\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)$					R_{Λ_c}				
	SM	SM	scalar LQ		vector LQ		SM	scalar LQ		vector LQ	
			P_A	P_C	R_A	R_B		P_A	P_C	R_A	R_B
1	2.50	0.74	0.95	0.93	0.95 ± 0.05	0.94 ± 0.05	0.30	0.38	0.37	0.38 ± 0.02	0.37 ± 0.02
2	2.67	0.73	0.93	0.91	0.93 ± 0.05	0.92 ± 0.05	0.27	0.35	0.34	0.35 ± 0.02	0.34 ± 0.02
3	5.16	1.39	1.77	1.73	1.78 ± 0.09	1.75 ± 0.09	0.27	0.35	0.34	0.35 ± 0.02	0.34 ± 0.02
4	5.74	1.50	1.92	1.88	1.93 ± 0.10	1.90 ± 0.10	0.26	0.34	0.33	0.34 ± 0.02	0.33 ± 0.02

Table 3. Predictions for the branching fractions (in unit of 10^{-2}) and the ratio R_{Λ_c} of $\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell$ ($\ell = e/\tau$) decays both within the SM and in the scalar/vector LQ scenarios, with the form factors taken from QCD sum rules [30].

For the $\Lambda_b \rightarrow \Lambda_c$ transition form factors, we firstly use the results obtained in QCD sum rules [30], together with the HQET relations among the form factors [31–33]. Four types of form-factor parametrizations for two values of the parameter κ , which is introduced to account for deviations from the factorization hypothesis for four-quark condensates, and for two choices of the continuum model are shown in table 2. For a comparison, we also adopt the latest lattice QCD results for the (axial-)vector form factors [26], where the q^2 dependence of the form factors is parameterized in a simplified z expansion [114], modified to account for pion-mass and lattice-spacing dependence. All relevant formulae and input data can be found in eq. (79) and tables VII–IX of ref. [26]. While for the (pseudo-)tensor form factors, since lattice result is unavailable so far, we still use the HQET relations to relate them to the corresponding (axial-)vector ones.

4.2 Numerical analyses

We now give our predictions for the branching fractions $\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell)$ and the ratio R_{Λ_c} both within the SM and in the scalar and the vector LQ scenarios in table 3 with the form factors taken from QCD sum rules [30], and in table 4 with the form factors taken from lattice QCD calculations [26]. The theoretical uncertainties in table 3 come only from the NP parameters given by eq. (2.9), whereas in table 4 we have also included the uncertainties from the form-factor parameters following the procedure recommended in [26]. Specifically, we have taken into account the correlation matrices between the form-factor parameters, and calculate the central values, statistical uncertainties, and total systematic uncertainties of any observable depending on these parameters, according to eqs. (82)–(84) specified in ref. [26].

From the numerical results given in tables 3 and 4, we can draw the following conclusions:

- The branching fractions are very sensitive to the form-factor parameterizations used in QCD sum rules [30]. The triangular region (with $\kappa = 3$ or $\kappa = 4$) for the continuum model gives more reliable predictions compared to the rectangular one, because, within the SM, the former leads to consistent results with that obtained using the lattice-based form factors, and also with the experimental data $\mathcal{B}(\Lambda_b \rightarrow \Lambda_c e \bar{\nu}_e) =$

$\mathcal{B}(\Lambda_b \rightarrow \Lambda_c e \bar{\nu}_e)$	SM		5.34 ± 0.33
$\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)$	SM		1.77 ± 0.09
	scalar LQ	P_A	2.26 ± 0.12
		P_C	2.22 ± 0.12
	vector LQ	R_A	2.27 ± 0.17
		R_B	2.24 ± 0.17
R_{Λ_c}	SM		0.33 ± 0.01
	scalar LQ	P_A	0.42 ± 0.01
		P_C	0.42 ± 0.01
	vector LQ	R_A	0.43 ± 0.02
		R_B	0.42 ± 0.02

Table 4. Same as in table 3 but with the form factors taken from lattice QCD calculation [26].

$(6.5^{+3.2}_{-2.5})\%$ [113]. Thus, from now on, we consider only the triangular continuum model with two different values of κ . The ratio R_{Λ_c} , on the other hand, is insensitive to the form-factor parameterizations, as is generally expected. It is also noted that the predicted R_{Λ_c} using the lattice-based form factors is a little bit larger than that obtained from QCD sum rules, both within the SM and in the two LQ scenarios.

- In the scalar LQ scenario, the predicted branching fraction $\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)$ is enhanced by about 28% (25%) in the P_A (P_C) case, no matter the form factors are taken from QCD sum rules or from the lattice QCD calculations. The slight difference between the two solutions results from the $\sim 1.2\%$ numerical difference in the dominated coefficient $|C_V^{\text{fit}}|$, which has been discussed in ref. [19]. As the decay modes with light leptons ($\ell = e, \mu$) are assumed to be free from the scalar LQ contribution, the ratio R_{Λ_c} is also enhanced by the corresponding percentages relative to the SM prediction.
- In the vector LQ scenario, compared to the SM prediction, the branching fraction $\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)$ is found to be enhanced by about 28% in the R_A and by about 27% in the R_B case, respectively. To understand this, we should note that there is only one ($V - A$) coupling in this scenario, and the resulting effective coefficients C'_V , corresponding to the two solutions R_A and R_B (cf. eq. (2.9)), are given, respectively, as

$$C'_V{}^{\text{fit}} = \begin{cases} 1.133 \pm 0.030, & \text{for } R_A \\ -1.124 \pm 0.030, & \text{for } R_B \end{cases} . \quad (4.1)$$

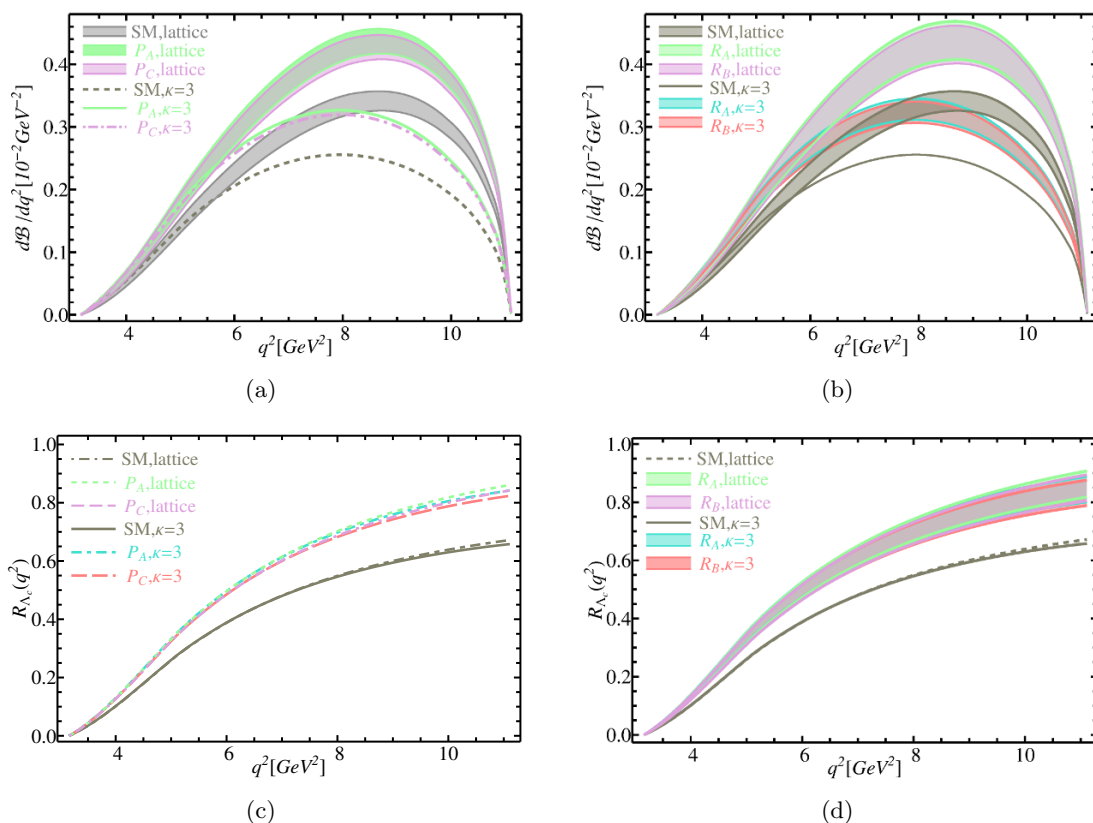


Figure 1. The q^2 distributions of the differential branching fraction $d\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)/dq^2$ ((a): in the scalar and (b): in the vector scenario) and the ratio $R_{\Lambda_c}(q^2)$ ((c): in the scalar scenario and (d): in the vector scenario). The bands in (a) and (c) due to the uncertainties of form-factor parameters obtained in lattice QCD, and in (b) and (d) also include the varyings of the NP parameters in the vector scenario

One can see clearly that, just like C_V^{fit} in scalar LQ scenario, $C_V^{\prime\text{fit}}$ also has nearly the same absolute values for the two solutions R_A and R_B , both enhancing the SM result by $\sim 13\%$, but the sign of solution R_B is flipped relative to the SM part.

- As the two effective couplings $|C_V^{\text{fit}}|$ and $|C_V^{\prime\text{fit}}|$ are both enhanced by about $12\% \sim 13\%$, compared to the SM part, they would give quite similar predictions for the other observables in $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ decay.

The q^2 dependences of the differential branching fraction $d\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)/dq^2$ and the ratio $R_{\Lambda_c}(q^2)$ are displayed in figure 1, both within the SM and in the two LQ scenarios. As the results based on the form-factor parametrizations with $\kappa = 3$ are similar to that with $\kappa = 4$, we show only the case with $\kappa = 3$. One can see that these two observables present the same features as the corresponding q^2 -integrated ones discussed above: the predicted $d\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)/dq^2$ using the lattice-based form factors are a little bit larger than that based on QCD sum rules, and are enhanced at $q^2 \sim 9 \text{ GeV}^2$ at most in both the two LQ scenarios. However, the ratio $R_{\Lambda_c}(q^2)$ is insensitive to the choices of the form factors.

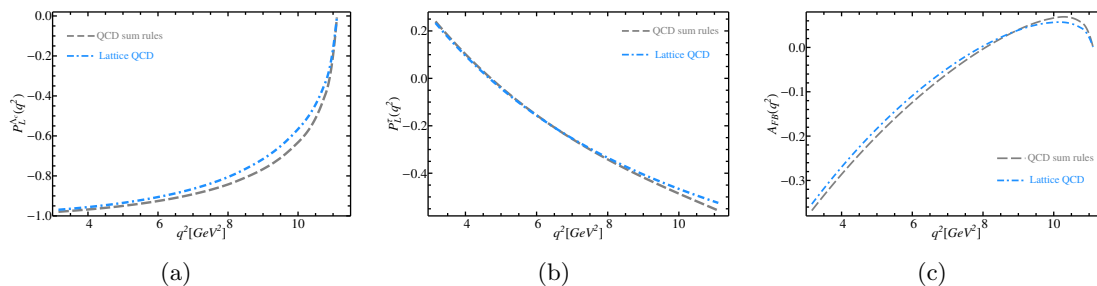


Figure 2. The q^2 dependences of the Λ_c (a) and τ (b) longitudinal polarizations as well as the lepton-side forward-backward asymmetry (c), both within the SM and in the two LQ scenarios.

Finally, we show in figure 2 the q^2 dependences of the Λ_c and τ longitudinal polarizations as well as the lepton-side forward-backward asymmetry. Since the resulting effective coefficients C_V and C'_V of the dominated $(V - A)$ couplings appear both in the numerator and in the denominator of these ratios, the NP effects are cancelled exclusively. At the same time, the form-factor dependences of these observables are reduced to a large extent, and all the four cases in QCD sum rules ($\kappa = 1, \dots, 4$) give almost the same curves for each observable, while being only slightly different from that obtained with the lattice-based form factors, as shown in figure 2. As a consequence, all these three observables are insensitive to the two LQ scenarios and behave nearly the same as in the SM.

5 Conclusions

As demonstrated in refs. [16–18], both the scalar and vector LQ scenarios could explain the anomalies observed in $\bar{B} \rightarrow D^{(*)}\tau\bar{\nu}_\tau$ and $\bar{B} \rightarrow \bar{K}\ell^+\ell^-$ decays and, for each scenario, there exist two best-fit solutions for the operator coefficients, because the other two solutions for scalar LQ are already excluded by the $B_c^- \rightarrow \tau^-\bar{\nu}_\tau$ decay [19]. To further explore these two interesting scenarios, in this paper, we have studied their effects in the semi-leptonic $\Lambda_b \rightarrow \Lambda_c\tau\bar{\nu}_\tau$ decay, which is induced by the same quark-level transition as in $\bar{B} \rightarrow D^{(*)}\tau\bar{\nu}_\tau$ decays. Besides the branching fraction $\mathcal{B}(\Lambda_b \rightarrow \Lambda_c\tau\bar{\nu}_\tau)$ and the ratio R_{Λ_c} , we have also discussed the q^2 distributions of these two observables, as well as the Λ_c and τ longitudinal polarizations and the lepton-side forward-backward asymmetry, using the $\Lambda_b \rightarrow \Lambda_c$ transition form factors from both the QCD sum rules and the latest lattice QCD calculations. In addition, we have also discussed the feasibility for the measurements of these observables at the LHC and the future e^+e^- colliders.

Using the best-fit solutions for the operator coefficients allowed by the current data of mesonic decays, we have found that the two LQ scenarios give the similar amounts of enhancements relative to the SM predictions for the branching fraction $\mathcal{B}(\Lambda_b \rightarrow \Lambda_c\tau\bar{\nu}_\tau)$ and the ratio R_{Λ_c} . The two best-fit solutions in each of the two scenarios are still found to be indistinguishable from each other based only on these two observables. On the other hand, both of the two LQ scenarios give nearly the same predictions as the SM for the Λ_c and τ longitudinal polarizations, as well as the lepton-side forward-backward asymmetry. As a consequence, we conclude that, while the two LQ scenarios could be distinguished

from the SM, it is quite difficult to distinguish between them using the semi-leptonic $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ decay.

With large numbers of Λ_b produced at the LHC and the future e^+e^- colliders, we expect that the two LQ scenarios could be further tested, and even differentiated from other NP explanations to the $R_{D^{(*)}}$ anomalies, with the measurement of $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ decay.

Acknowledgments

We thank Prof. Yuehong Xie and Dr. Jiesheng Yu for useful discussions about the measurements of the $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ decay at the LHC and future e^+e^- colliders. The work is supported by the National Natural Science Foundation of China (NSFC) under contract Nos. 11675061, 11435003, 11225523 and 11521064. XL acknowledges the hospitality of the Munich Institute for Astro- and Particle Physics (MIAPP) of the DFG cluster of excellence ‘‘Origin and Structure of the Universe’’, where this work was finalized.

A $\Lambda_b \rightarrow \Lambda_c$ transition form factors

The hadronic matrix elements of the vector and axial-vector currents between the two spin-half baryons Λ_b and Λ_c can be parameterized in terms of three form factors, respectively, as [24]

$$\langle \Lambda_c, \lambda_2 | \bar{c} \gamma_\mu b | \Lambda_b, \lambda_1 \rangle = \bar{u}_2(p_2, \lambda_2) \left[F_1^V(q^2) \gamma_\mu - \frac{F_2^V(q^2)}{M_{\Lambda_b}} i \sigma_{\mu\nu} q^\nu + \frac{F_3^V(q^2)}{M_{\Lambda_b}} q_\mu \right] u_1(p_1, \lambda_1), \quad (\text{A.1})$$

$$\langle \Lambda_c, \lambda_2 | \bar{c} \gamma_\mu \gamma_5 b | \Lambda_b, \lambda_1 \rangle = \bar{u}_2(p_2, \lambda_2) \left[F_1^A(q^2) \gamma_\mu - \frac{F_2^A(q^2)}{M_{\Lambda_b}} i \sigma_{\mu\nu} q^\nu + \frac{F_3^A(q^2)}{M_{\Lambda_b}} q_\mu \right] \gamma_5 u_1(p_1, \lambda_1), \quad (\text{A.2})$$

where $\sigma_{\mu\nu} = \frac{i}{2}(\gamma_\mu \gamma_\nu - \gamma_\nu \gamma_\mu)$, $q = p_1 - p_2$ is the four-momentum transfer, and $\lambda_i = \pm \frac{1}{2}$ ($i = 1, 2$) denote the helicities of the Λ_b and Λ_c baryons, respectively. Using the equations of motion, we can then obtain the hadronic matrix elements of the scalar and pseudo-scalar currents between the two baryons, which are given, respectively, as

$$\langle \Lambda_c, \lambda_2 | \bar{c} b | \Lambda_b, \lambda_1 \rangle = \frac{1}{m_b - m_c} \bar{u}_2(p_2, \lambda_2) \left[F_1^V(q^2)(M_1 - M_2) + \frac{F_3^V(q^2)}{M_{\Lambda_b}} q^2 \right] u_1(p_1, \lambda_1), \quad (\text{A.3})$$

$$\langle \Lambda_c, \lambda_2 | \bar{c} \gamma_5 b | \Lambda_b, \lambda_1 \rangle = \frac{1}{m_b + m_c} \bar{u}_2(p_2, \lambda_2) \left[F_1^A(q^2)(M_1 + M_2) - \frac{F_3^A(q^2)}{M_{\Lambda_b}} q^2 \right] \gamma_5 u_1(p_1, \lambda_1), \quad (\text{A.4})$$

where m_b and m_c are the current quark masses evaluated at the scale $\mu \sim m_b$.

The hadronic matrix elements of the tensor and pseudo-tensor currents between the Λ_b and Λ_c baryons can be generally parameterized as [115]

$$\begin{aligned} \langle \Lambda_c, \lambda_2 | \bar{c} i \sigma_{\mu\nu} b | \Lambda_b, \lambda_1 \rangle &= \bar{u}_2(p_2, \lambda_2) \left[f_T^V i \sigma_{\mu\nu} + f_T^V (\gamma_\mu q_\nu - \gamma_\nu q_\mu) \right. \\ &\quad \left. + f_T^S (P_\mu q_\nu - P_\nu q_\mu) \right] u_1(p_1, \lambda_1), \end{aligned} \quad (\text{A.5})$$

$$\begin{aligned} \langle \Lambda_c, \lambda_2 | \bar{c} i \sigma_{\mu\nu} \gamma_5 b | \Lambda_b, \lambda_1 \rangle &= \bar{u}_2(p_2, \lambda_2) [g_T i \sigma_{\mu\nu} + g_T^V (\gamma_\mu q_\nu - \gamma_\nu q_\mu) \\ &\quad + g_T^S (P_\mu q_\nu - P_\nu q_\mu)] \gamma_5 u_1(p_1, \lambda_1), \end{aligned} \quad (\text{A.6})$$

where $P = p_1 + p_2$. The $\Lambda_b \rightarrow \Lambda_c$ transition form factors have also been studied based on the HQET [31–33], and the following relations among the form factors can be found [115]:

$$F_1^V = F_1^A = f_T = g_T, \quad (\text{A.7})$$

$$F_2^V = F_2^A = -F_3^V = -F_3^A, \quad (\text{A.8})$$

$$f_T^V = g_T^V = f_T^S = g_T^S = 0. \quad (\text{A.9})$$

An alternate helicity-based definition of the $\Lambda_b \rightarrow \Lambda_c$ form factors can be found in ref. [116], and the explicit relations between these two sets of form factors are also given in ref. [116].

In this paper, we use the results obtained both in the QCD sum rules [116] and in the most recent lattice QCD calculation with $2 + 1$ dynamical flavours [26]. However, since there are currently no lattice results for the (pseudo-)tensor form factors yet, we still use the following HQET relations to relate them to the corresponding (axial-)vector ones,

$$f_T = g_T = F_1^V = \frac{(M_{\Lambda_b} + M_{\Lambda_c})^2 f_+ - q^2 f_\perp}{(M_{\Lambda_b} + M_{\Lambda_c})^2 - q^2}, \quad f_T^V = g_T^V = f_T^S = g_T^S = 0. \quad (\text{A.10})$$

B Helicity-dependent differential decay rates

In order to discuss the Λ_c and τ polarizations, we need the helicity-dependent differential decay rates, which are collected below (normalized by the prefactor N defined in eq. (3.7)):

$$\begin{aligned} \frac{d\Gamma^{\lambda_2=1/2}}{dq^2} &= \frac{m_\ell^2}{q^2} \left[\frac{4}{3} C_V^2 \left(H_{\frac{1}{2},+}^2 + H_{\frac{1}{2},0}^2 + 3H_{\frac{1}{2},t}^2 \right) + \frac{2}{3} C_T^2 \left(H_{\frac{1}{2},+,-}^{T2} + H_{\frac{1}{2},0,t}^{T2} + H_{\frac{1}{2},+,0}^{T2} \right. \right. \\ &\quad \left. \left. + H_{\frac{1}{2},+,t}^{T2} + 2H_{\frac{1}{2},+,-}^T - H_{\frac{1}{2},0,t}^T + 2H_{\frac{1}{2},+,0}^T H_{\frac{1}{2},+,t}^T \right) \right] \\ &\quad + \frac{8}{3} C_V^2 \left(H_{\frac{1}{2},0}^2 + H_{\frac{1}{2},+}^2 \right) + 4C_S^2 H_{\frac{1}{2},0}^{SP2} \\ &\quad + \frac{C_T^2}{3} \left(H_{\frac{1}{2},+,-}^{T2} + H_{\frac{1}{2},0,t}^{T2} + H_{\frac{1}{2},+,0}^{T2} + H_{\frac{1}{2},+,t}^{T2} + 2H_{\frac{1}{2},+,-}^T - H_{\frac{1}{2},0,t}^T + 2H_{\frac{1}{2},+,0}^T H_{\frac{1}{2},+,t}^T \right) \\ &\quad + \frac{4m_\ell}{\sqrt{q^2}} \left[C_V C_T \left(H_{\frac{1}{2},0} H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0} H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+} H_{\frac{1}{2},+,0}^T + H_{\frac{1}{2},+} H_{\frac{1}{2},+,t}^T \right) \right. \\ &\quad \left. + 2C_V C_S \left(H_{\frac{1}{2},t} H_{\frac{1}{2},0}^{SP} \right) \right], \\ \frac{d\Gamma^{\lambda_2=-1/2}}{dq^2} &= \frac{m_\ell^2}{q^2} \left[\frac{4}{3} C_V^2 \left(H_{-\frac{1}{2},-}^2 + H_{-\frac{1}{2},0}^2 + 3H_{-\frac{1}{2},t}^2 \right) + \frac{2}{3} C_T^2 \left(H_{-\frac{1}{2},+,-}^{T2} + H_{-\frac{1}{2},0,t}^{T2} + H_{-\frac{1}{2},0,-}^{T2} \right. \right. \\ &\quad \left. \left. + H_{-\frac{1}{2},-,t}^{T2} + 2H_{-\frac{1}{2},+,-}^T - H_{-\frac{1}{2},0,t}^T + 2H_{-\frac{1}{2},0,-}^T - H_{-\frac{1}{2},-,t}^T \right) \right] \\ &\quad + \frac{8}{3} C_V^2 \left(H_{-\frac{1}{2},-}^2 + H_{-\frac{1}{2},0}^2 \right) + 4C_S^2 H_{-\frac{1}{2},0}^{SP2} \\ &\quad + \frac{C_T^2}{3} \left(H_{-\frac{1}{2},+,-}^{T2} + H_{-\frac{1}{2},0,t}^{T2} + H_{-\frac{1}{2},0,-}^{T2} + H_{-\frac{1}{2},-,t}^{T2} \right. \\ &\quad \left. + 2H_{-\frac{1}{2},+,-}^T - H_{-\frac{1}{2},0,t}^T + 2H_{-\frac{1}{2},0,-}^T - H_{-\frac{1}{2},-,t}^T \right) \end{aligned}$$

$$\begin{aligned}
 & + \frac{4m_\ell}{\sqrt{q^2}} \left[C_V C_T \left(H_{-\frac{1}{2},0} H_{-\frac{1}{2},+,-}^T + H_{-\frac{1}{2},0} H_{-\frac{1}{2},0,t}^T \right. \right. \\
 & \quad \left. \left. + H_{-\frac{1}{2},-} H_{-\frac{1}{2},0,-}^T + H_{-\frac{1}{2},-} H_{-\frac{1}{2},-,t}^T \right) + 2C_V C_S \left(H_{-\frac{1}{2},t} H_{-\frac{1}{2},0}^{SP} \right) \right], \\
 \frac{d\Gamma^{\lambda_\tau=1/2}}{dq^2} &= \frac{m_\ell^2}{q^2} C_V^2 \left[\frac{4}{3} \left(H_{\frac{1}{2},+}^2 + H_{\frac{1}{2},0}^2 + H_{-\frac{1}{2},-}^2 + H_{-\frac{1}{2},0}^2 \right) + 4 \left(H_{\frac{1}{2},t}^2 + H_{-\frac{1}{2},t}^2 \right) \right] \\
 & + 4C_S^2 \left(H_{\frac{1}{2},0}^{SP2} + H_{-\frac{1}{2},0}^{SP2} \right) \\
 & + \frac{C_T^2}{3} \left[H_{\frac{1}{2},+,-}^{T2} + H_{\frac{1}{2},0,t}^{T2} + H_{\frac{1}{2},+,0}^{T2} + H_{\frac{1}{2},+,t}^{T2} + H_{-\frac{1}{2},+,-}^{T2} + H_{-\frac{1}{2},0,t}^{T2} \right. \\
 & \quad \left. + H_{-\frac{1}{2},0,-}^{T2} + H_{-\frac{1}{2},-,t}^{T2} + 2 \left(H_{\frac{1}{2},+,-}^T H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+,0}^T H_{\frac{1}{2},+,t}^T \right. \right. \\
 & \quad \left. \left. + H_{-\frac{1}{2},+,-}^T H_{-\frac{1}{2},0,t}^T + H_{-\frac{1}{2},0,-}^T H_{-\frac{1}{2},-,t}^T \right) \right] \\
 & + \frac{4m_\ell}{3\sqrt{q^2}} \left[6C_V C_S \left(H_{\frac{1}{2},t} H_{\frac{1}{2},0}^{SP} + H_{-\frac{1}{2},t} H_{-\frac{1}{2},0}^{SP} \right) + C_V C_T \left(H_{\frac{1}{2},0} H_{\frac{1}{2},+,-}^T \right. \right. \\
 & \quad \left. \left. + H_{\frac{1}{2},0} H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+} H_{\frac{1}{2},+,0}^T + H_{\frac{1}{2},+} H_{\frac{1}{2},+,t}^T + H_{-\frac{1}{2},0} H_{-\frac{1}{2},+,-}^T \right. \right. \\
 & \quad \left. \left. + H_{-\frac{1}{2},0} H_{-\frac{1}{2},0,t}^T + H_{-\frac{1}{2},-} H_{-\frac{1}{2},0,-}^T + H_{-\frac{1}{2},-} H_{-\frac{1}{2},-,t}^T \right) \right], \\
 \frac{d\Gamma^{\lambda_\tau=-1/2}}{dq^2} &= \frac{8C_V^2}{3} \left(H_{\frac{1}{2},+}^2 + H_{\frac{1}{2},0}^2 + H_{-\frac{1}{2},-}^2 + H_{-\frac{1}{2},0}^2 \right) \\
 & + \frac{2m_\ell^2}{3q^2} C_T^2 \left[H_{\frac{1}{2},+,-}^{T2} + H_{\frac{1}{2},0,t}^{T2} + H_{\frac{1}{2},+,0}^{T2} + H_{\frac{1}{2},+,t}^{T2} \right. \\
 & \quad \left. + H_{-\frac{1}{2},+,-}^{T2} + H_{-\frac{1}{2},0,t}^{T2} + H_{-\frac{1}{2},0,-}^{T2} + H_{-\frac{1}{2},-,t}^{T2} + 2 \left(H_{\frac{1}{2},+,-}^T H_{\frac{1}{2},0,t}^T \right. \right. \\
 & \quad \left. \left. + H_{\frac{1}{2},+,0}^T H_{\frac{1}{2},+,t}^T + H_{-\frac{1}{2},+,-}^T H_{-\frac{1}{2},0,t}^T + H_{-\frac{1}{2},0,-}^T H_{-\frac{1}{2},-,t}^T \right) \right] \\
 & + \frac{8m_\ell}{3\sqrt{q^2}} C_V C_T \left(H_{\frac{1}{2},0} H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0} H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+} H_{\frac{1}{2},+,0}^T + H_{\frac{1}{2},+} H_{\frac{1}{2},+,t}^T \right. \\
 & \quad \left. + H_{-\frac{1}{2},0} H_{-\frac{1}{2},+,-}^T + H_{-\frac{1}{2},0} H_{-\frac{1}{2},0,t}^T + H_{-\frac{1}{2},-} H_{-\frac{1}{2},0,-}^T + H_{-\frac{1}{2},-} H_{-\frac{1}{2},-,t}^T \right). \tag{B.1}
 \end{aligned}$$

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