Investigation of effects of new physics in $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ decay

Xiao-Long Mu¹, Ying Li^{1,2,*}, Zhi-Tian Zou¹ and Bin Zhu¹

¹Department of Physics, Yantai University, Yantai 264005, China

²Center for High Energy Physics, Peking University, Beijing 100871, China

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Recent experimental results of $\mathcal{R}(D^{(*)})$ deviate from the standard model (SM) by 3.1σ , suggesting a new physics (NP) that affects the $b \rightarrow c \tau \bar{\nu}_\tau$ transition. Motivated by this, we investigate the possible NP effects in the $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ decay. For this purpose, assuming the neutrinos are left-handed, we calculate in detail the helicity amplitudes of $\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell$ ($\ell = e^-, \mu^-, \text{ and } \tau^-$) decays with all possible four-fermion operators. Within the latest results of $\Lambda_b \rightarrow \Lambda_c$ form factors from lattice QCD calculations, we study these decays in a model-independent manner. The differential and total branching fractions, the longitudinal polarizations of final leptons and hadrons, the forward-backward asymmetries in the lepton side, the convexity parameters, and the ratio $\mathcal{R}(\Lambda_c)$ are calculated. In SM, we obtain the ratio $\mathcal{R}(\Lambda_c) = 0.33 \pm 0.01$. Supposing that NP only affects the third generation fermions, we present the correlations among $\mathcal{R}(D)$, $\mathcal{R}(D^*)$, and $\mathcal{R}(\Lambda_c)$, as the $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ and $B \rightarrow D^{(*)} \tau \bar{\nu}_\tau$ are all induced by $b \rightarrow c \ell \bar{\nu}_\ell$. We perform a minimum χ^2 fit of the Wilson coefficient of each operator to the latest experimental data of different observables, including the ratios $\mathcal{R}(D^{(*)})$ and $\mathcal{R}(J/\psi)$ and the τ polarization $P_\tau(D^*)$. It is found that the left-handed scalar operator $\mathcal{O}_{SL}$ affects the branching fraction remarkably, and the ratio $\mathcal{R}(\Lambda_c)$ can be enhanced by 30%. For other operators, the ratio amounts to 0.38 ± 0.02 , which is larger than prediction of SM by 20%. Using the fitted values of the Wilson coefficients of the single NP operators, we also give a prognosis for the physical observables of $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$, including the ratio $\mathcal{R}(\Lambda_c)$, forward-backward asymmetry, and other polarized observables as well as the differential branching fraction. Furthermore, we also study the effects of three typical NP models on the ratio and the differential branching fraction of $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$. We hope our results can be tested in the current LHCb experiment and the future high energy experiments.

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I. INTRODUCTION

After the discovery of the Higgs boson, one of the most important tasks is searching for the possible new physics (NP) beyond the standard model (SM). So far, at the LHC, no one has directly observed any signals of new particles beyond the SM. However, there are currently a number of anomalies in semileptonic b -flavor hadron decays, which have attracted great attention in the scientific community. The unexpected deviations seem to appear in both $b \rightarrow c$ and $b \rightarrow s$ semileptonic decay transitions when different generations of leptons are involved; see Refs. [1–3] for recent reviews. Unlike the flavor-changing neutral-current $b \rightarrow s$ transition that has been thought of as an ideal plate for probing the effect of NP, a $b \rightarrow c$ transition is a tree

level process by the exchange of a W boson in SM. Primordially, in order to study the charged Higgs contribution, the ratios are defined as [4,5]

$$\mathcal{R}(D^{(*)}) = \frac{\mathcal{B}(B \rightarrow D^{(*)} \tau \bar{\nu})}{\mathcal{B}(B \rightarrow D^{(*)} \ell \bar{\nu})}, \quad \ell = e, \mu. \quad (1)$$

For these two ratios, the uncertainties from the Cabibbo-Kobayashi-Maskawa (CKM) matrix element V_{cb} can be removed, and ones from the hadronic transition form factors can also be largely reduced. In SM, based on the heavy-quark effective theory or the lattice QCD approach, they have been studied explicitly [6–11], and the theoretical averaged results are given by [12]

$$\mathcal{R}(D)^{\text{SM}} = 0.299^{+0.003}_{-0.003}, \quad \text{and} \quad \mathcal{R}(D^*)^{\text{SM}} = 0.258^{+0.005}_{-0.005}. \quad (2)$$

However, from 2012, measurements from BABAR [13,14], Belle [15–17], and LHCb [18,19] showed that there are deviations between experimental data and predictions of

*liying@ytu.edu.cn

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SM, and the combination results performed by the Heavy Flavor Averaging Group (HFLAV) [12] are

$$\begin{aligned}\mathcal{R}(D)^{\text{avg}} &= 0.407 \pm 0.039 \pm 0.024 \\ \text{and } \mathcal{R}(D^*)^{\text{avg}} &= 0.306 \pm 0.013 \pm 0.007,\end{aligned}\quad (3)$$

with a correlation of -20% , which shows a tension of 3.8σ with SM predictions. Very recently, at Moriond EW 2019, the Belle Collaboration released the latest measurement of $\mathcal{R}(D)$ and $\mathcal{R}(D^*)$ [20],

$$\begin{aligned}\mathcal{R}(D)^{\text{Belle,2019}} &= 0.307 \pm 0.037 \pm 0.016 \\ \text{and } \mathcal{R}(D^*)^{\text{Belle,2019}} &= 0.283 \pm 0.018 \pm 0.014,\end{aligned}\quad (4)$$

with a correlation of -53% , and the combined results compatible with SM at the 1.2σ level. Including these new measurements in the global average leads to [12]

$$\begin{aligned}\mathcal{R}(D)^{\text{avg,new}} &= 0.340 \pm 0.027 \pm 0.013 \\ \text{and } \mathcal{R}(D^*)^{\text{avg,new}} &= 0.299 \pm 0.013 \pm 0.008,\end{aligned}\quad (5)$$

and the significance of the anomaly amounts to 3.1σ relative to the above SM prediction. In addition, LHCb also reported a value of the ratio as [21]

$$\mathcal{R}(J/\psi) = \frac{\mathcal{B}(B_c \rightarrow J/\psi \tau \bar{\nu})}{\mathcal{B}(B_c \rightarrow J/\psi \ell \bar{\nu})} = 0.71 \pm 0.17 \pm 0.18. \quad (6)$$

This result deviates 2σ away from the SM predictions, which lies in the range of 0.23 to 0.28, where the uncertainties arise from the choice of the modeling approach for the $B_c \rightarrow J/\psi$ form factors [22–27].

Although these deviations are perhaps from the uncertainties of hadronic transition form factors in SM, they might imply that the lepton flavor universality (LFU) is violated, which is the hint of the existence of NP, because the LFU is one of the major characters of SM. Due to the fact that there is no any similar discrepancy in K and π semileptonic and purely leptonic decays, or in electroweak precision observables, most of us believe that the LFU violation only appears in the third generation fermions. In this context, there are numerous works that had been carried out based on model-independence approaches [9,26,28–50] or special models by introducing new particles such as a charged Higgs boson [5,51–55], leptiquarks [56–79], or new vector bosons [80–88].

If the $\mathcal{R}(D^{(*)})$ anomalies are due to NP, it is natural for us to ask whether its effects can be shown up in other $b \rightarrow c \tau \bar{\nu}_\tau$ transition decays, such as $B_c \rightarrow \eta_c \tau \bar{\nu}_\tau$ [25,26,28,89,90], $B_s \rightarrow D_s^{(*)} \tau \bar{\nu}_\tau$ [91–93], and $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ [94–106]. In this work, we shall study the $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ comprehensively in model-dependence and independence manners, as well as $\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell$ ($\ell = e, \mu$). In this work, we will calculate the total rates, the differential decay distributions, the

longitudinal polarizations of the final states Λ_c and τ , and the forward-backward asymmetries in the lepton side, with and without contributions of NP. Of course, the correlations between $\mathcal{R}(D^{(*)})$ and $\mathcal{R}(\Lambda_b)$ will be given. Compared with previous studies, there are some improvements in the current work: (i) we calculate the complete amplitude including the contributions of scalar and tensor operators, and the latter was usually neglected in many previous studies; (ii) we present the relations between different conventions of form factors, which are nontrivial in the calculations; (iii) the $\Lambda_b \rightarrow \Lambda_c$ form factors we adopted are the latest results from lattice QCD calculations; (iv) in discussing the effects of NP, we refit each Wilson coefficient individually within the latest experimental averaged results, and the form factors of $B \rightarrow D^{(*)}$ in fitting are from the heavy quark effective theory including the next-to-leading power $\mathcal{O}(1/m_b)$ and next-to-leading order $\mathcal{O}(\alpha_s)$ corrections.

Our paper is organized as follows: In Sec. II, the analytical formulas in this work are presented, such as the effective Hamiltonian, the form factors, and the helicity amplitudes. All definitions of physical observables are also given in this section. In Sec. III, we will present the numerical results and discussions. At last, we will draw conclusions in Sec. IV.

II. ANALYTICAL FORMULAS

A. The effective Hamiltonian

Within the presence of NP, the effective Lagrangian can be extended by incorporating new operator basis that includes all possible four-fermion interactions. Due to the absence of experimental evidence of deviations from the SM in tree-level transitions involving light leptons such as the precision measurements of the ratio $\mathcal{B}(\tau \rightarrow \mu \nu_\tau \bar{\nu}_\mu)/\mathcal{B}(\tau \rightarrow e \nu_\tau \bar{\nu}_e) = 0.9762 \pm 0.0028$ [107], the new physics is generally supposed to appear in the third generation fermions. Therefore, without considering the effects from the right-handed neutrinos, the effective Lagrangian of $b \rightarrow c \tau \bar{\nu}_\tau$ can be written as

$$\begin{aligned}\mathcal{L}_{\text{eff}} &= -\frac{4G_F}{\sqrt{2}} V_{cb} [(1 + C_{VL})\mathcal{O}_{VL} + C_{VR}\mathcal{O}_{VR} + C_{SL}\mathcal{O}_{SL} \\ &\quad + C_{SR}\mathcal{O}_{SR} + C_T\mathcal{O}_T] + \text{H.c.},\end{aligned}\quad (7)$$

where G_F is the Fermi constant and V_{cb} is the CKM matrix element. The four-fermion operators can be defined as

$$\begin{aligned}\mathcal{O}_{VL} &= (\bar{c}\gamma^\mu P_L b)(\bar{\tau}\gamma_\mu P_L \nu_\tau), & \mathcal{O}_{VR} &= (\bar{c}\gamma^\mu P_R b)(\bar{\tau}\gamma_\mu P_L \nu_\tau), \\ \mathcal{O}_{SL} &= (\bar{c}P_L b)(\bar{\tau}P_L \nu_\tau), & \mathcal{O}_{SR} &= (\bar{c}P_R b)(\bar{\tau}P_L \nu_\tau), \\ \mathcal{O}_T &= (\bar{c}\sigma^{\mu\nu} P_L b)(\bar{\tau}\sigma_{\mu\nu} P_L \nu_\tau),\end{aligned}\quad (8)$$

where $P_{L,R} = (1 \mp \gamma_5)/2$ and C_i ($i = VL, VR, SL, SR$, and T) are the corresponding Wilson coefficients at the

scale $\mu = m_b$, with $C_i = 0$ in SM. It should be stressed that recent studies [108–110] based on the effective field theory of SM (SMEFT) [111,112] showed that the operator $\mathcal{O}_{VR}$ does not contribute to the LFU violation at leading order; however, we here include this operator for completeness.

B. Form factors

In the calculation, the most important inputs are the hadronic 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 parametrized in terms of three form factors, respectively, as [94]

$$\langle \Lambda_c, \lambda_2 | \bar{c} \gamma_\mu b | \Lambda_b, \lambda_1 \rangle = \bar{u}_2(p_2, \lambda_2) \left[F_1^V(q^2) \gamma_\mu - \frac{F_2^V(q^2)}{M_{\Lambda_b}} i \sigma_{\mu\nu} q^\nu + \frac{F_3^V(q^2)}{M_{\Lambda_b}} q_\mu \right] u_1(p_1, \lambda_1), \quad (9)$$

$$\langle \Lambda_c, \lambda_2 | \bar{c} \gamma_\mu \gamma_5 b | \Lambda_b, \lambda_1 \rangle = \bar{u}_2(p_2, \lambda_2) \left[F_1^A(q^2) \gamma_\mu - \frac{F_2^A(q^2)}{M_{\Lambda_b}} i \sigma_{\mu\nu} q^\nu + \frac{F_3^A(q^2)}{M_{\Lambda_b}} q_\mu \right] \gamma_5 u_1(p_1, \lambda_1), \quad (10)$$

where $\sigma_{\mu\nu} = \frac{i}{2}(\gamma_\mu \gamma_\nu - \gamma_\nu \gamma_\mu)$, $q = p_1 - p_2$. $\lambda_i = \pm \frac{1}{2}$ ($i = 1, 2$) denote the helicities of the Λ_b and Λ_c baryons, respectively. Using the equations of motion, the ones of the scalar and pseudoscalar currents can be obtained as

$$\langle \Lambda_c, \lambda_2 | \bar{c} b | \Lambda_b, \lambda_1 \rangle = \frac{1}{m_b - m_c} \bar{u}_2(p_2, \lambda_2) \left[F_1^V(q^2) M_- + \frac{F_3^V(q^2)}{M_{\Lambda_b}} q^2 \right] u_1(p_1, \lambda_1), \quad (11)$$

$$\langle \Lambda_c, \lambda_2 | \bar{c} \gamma_5 b | \Lambda_b, \lambda_1 \rangle = \frac{1}{m_b + m_c} \bar{u}_2(p_2, \lambda_2) \left[F_1^A(q^2) M_+ - \frac{F_3^A(q^2)}{M_{\Lambda_b}} q^2 \right] \gamma_5 u_1(p_1, \lambda_1), \quad (12)$$

and m_b and m_c are the current quark masses evaluated at the scale $\mu \sim m_b$, and $M_\pm = M_{\Lambda_b} \pm M_{\Lambda_c}$. Similarly, the hadronic matrix elements of the tensor and pseudotensor currents between the Λ_b and Λ_c baryons can be generally parametrized into four form factors as [95]

$$\begin{aligned} \langle \Lambda_c, \lambda_2 | \bar{c} i \sigma_{\mu\nu} b | \Lambda_b, \lambda_1 \rangle &= \bar{u}_2(p_2, \lambda_2) [F_T i \sigma_{\mu\nu} + F_T^V(\gamma_\mu q_\nu - \gamma_\nu q_\mu) + F_T^P(\gamma_\mu P_\nu - \gamma_\nu P_\mu) \\ &\quad + F_T^S(P_\mu q_\nu - P_\nu q_\mu)] u_1(p_1, \lambda_1), \end{aligned} \quad (13)$$

$$\begin{aligned} \langle \Lambda_c, \lambda_2 | \bar{c} i \sigma_{\mu\nu} \gamma_5 b | \Lambda_b, \lambda_1 \rangle &= \bar{u}_2(p_2, \lambda_2) [G_T i \sigma_{\mu\nu} + G_T^V(\gamma_\mu q_\nu - \gamma_\nu q_\mu) + G_T^P(\gamma_\mu P_\nu - \gamma_\nu P_\mu) \\ &\quad + G_T^S(P_\mu q_\nu - P_\nu q_\mu)] \gamma_5 u_1(p_1, \lambda_1), \end{aligned} \quad (14)$$

where $P = p_1 + p_2$. Noted that the terms F_T^P and G_T^P had been missed in some literatures, such as in Ref. [95]. Furthermore, the F_T^i and G_T^i are related by using the identity $\sigma^{\mu\nu} \gamma_5 = -\frac{i}{2} \epsilon^{\mu\nu\alpha\beta} \sigma_{\alpha\beta}$, and the relations are given as

$$G_T = F_T - F_T^V M_+ - F_T^P M_- - F_T^S Q_+, \quad G_T^V = F_T^P - F_T^S M_-, \quad G_T^P = F_T^V + F_T^S M_+, \quad G_T^S = -F_T^S, \quad (15)$$

where $Q_\pm = M_\pm^2 - q^2$, which means that only four of them are linearly independent.

Alternatively, one can use the helicity-based definition of the $\Lambda_b \rightarrow \Lambda_c$ form factors [113], then the matrix elements of the vector and axial vector currents can be written in terms of another six helicity form factors f_+ , $f_\perp$, f_0 , g_+ , $g_\perp$, and g_0 as follows:

$$\begin{aligned} \langle \Lambda_c, \lambda_2 | \bar{c} \gamma_\mu b | \Lambda_b, \lambda_1 \rangle &= \bar{u}_2(p_2, \lambda_2) \left[f_0(q^2) M_- \frac{q^\mu}{q^2} + f_+(q^2) \frac{M_+}{Q_+} \left(P^\mu - M_+ M_- \frac{q^\mu}{q^2} \right) \right. \\ &\quad \left. + f_\perp(q^2) \left(\gamma^\mu - \frac{2M_{\Lambda_c}}{Q_+} p_1^\mu - \frac{2M_{\Lambda_b}}{Q_+} p_2^\mu \right) \right] \bar{u}_1(p_1, \lambda_1), \end{aligned} \quad (16)$$

$$\begin{aligned} \langle \Lambda_c, \lambda_2 | \bar{c} \gamma^\mu \gamma_5 b | \Lambda_b, \lambda_1 \rangle &= -\bar{u}_2(p_2, \lambda_2) \gamma_5 \left[g_0(q^2) M_+ \frac{q^\mu}{q^2} + g_+(q^2) \frac{M_-}{Q_-} \left(P^\mu - M_+ M_- \frac{q^\mu}{q^2} \right) \right. \\ &\quad \left. + g_\perp(q^2) \left(\gamma^\mu + \frac{2M_{\Lambda_c}}{Q_-} p_1^\mu - \frac{2M_{\Lambda_b}}{Q_-} p_2^\mu \right) \right] \bar{u}_1(p_1, \lambda_1). \end{aligned} \quad (17)$$

Again, within the equation of motion, the matrix elements of the scalar and pseudoscalar currents can be obtained as

$$\langle \Lambda_c, \lambda_2 | \bar{c} b | \Lambda_b, \lambda_1 \rangle = f_0(q^2) \frac{M_-}{m_b - m_c} \bar{u}_2(p_2, \lambda_2) \bar{u}_1(p_1, \lambda_1), \quad (18)$$

$$\langle \Lambda_c, \lambda_2 | \bar{c} \gamma_5 b | \Lambda_b, \lambda_1 \rangle = g_0(q^2) \frac{M_+}{m_b + m_c} \bar{u}_2(p_2, \lambda_2) \gamma_5 \bar{u}_1(p_1, \lambda_1). \quad (19)$$

Similarly, the matrix elements of the tensor current can also be written within the four helicity form factors h_+ , $h_\perp$, $\tilde{h}_+$, $\tilde{h}_\perp$ [98],

$$\begin{aligned} \langle \Lambda_c, \lambda_2 | \bar{c} i \sigma_{\mu\nu} b | \Lambda_b, \lambda_1 \rangle = & \bar{u}_2(p_2, \lambda_2) \left[2h_+(q^2) \frac{p_1^\mu p_2^\nu - p_1^\nu p_2^\mu}{Q_+} \right. \\ & + h_\perp(q^2) \left(\frac{M_+}{q^2} (q^\mu \gamma^\nu - q^\nu \gamma^\mu) - 2 \left(\frac{1}{q^2} + \frac{1}{Q_+} \right) (p_1^\mu p_2^\nu - p_1^\nu p_2^\mu) \right) \\ & + \tilde{h}_+(q^2) \left(i \sigma^{\mu\nu} - \frac{2}{Q_-} (M_{\Lambda_b} (p_2^\mu \gamma^\nu - p_2^\nu \gamma^\mu) - M_{\Lambda_c} (p_1^\mu \gamma^\nu - p_1^\nu \gamma^\mu) + p_1^\mu p_2^\nu - p_1^\nu p_2^\mu) \right) \\ & + \tilde{h}_\perp(q^2) \frac{M_-}{q^2 Q_-} ((M_+ M_- - q^2) (\gamma^\mu p_1^\nu - \gamma^\nu p_1^\mu) \\ & \left. - (M_+ M_- + q^2) (\gamma^\mu p_2^\nu - \gamma^\nu p_2^\mu) + 2M_- (p_1^\mu p_2^\nu - p_1^\nu p_2^\mu) \right) \bar{u}_1(p_1, \lambda_1). \end{aligned} \quad (20)$$

Of course, these two definitions are physically equivalent, and the relations between two groups of form factors are given as

$$\begin{aligned} F_1^V(q^2) &= \frac{f_+(q^2) M_+^2 - f_\perp(q^2) q^2}{Q_+}, & F_2^V(q^2) &= \frac{(f_\perp(q^2) - f_+(q^2)) M_{\Lambda_b} M_+}{Q_+}, \\ F_3^V(q^2) &= \frac{M_- M_{\Lambda_b} [(f_0(q^2) - f_+(q^2)) M_+^2 - (f_0(q^2) - f_\perp(q^2)) q^2]}{q^2 Q_+}, & F_1^A(q^2) &= \frac{g_+(q^2) M_-^2 - g_\perp(q^2) q^2}{Q_-}, \\ F_2^A(q^2) &= \frac{(g_\perp(q^2) - g_+(q^2)) M_{\Lambda_b} M_-}{Q_-}, & F_3^A(q^2) &= \frac{M_+ M_{\Lambda_b} [(g_+(q^2) - g_0(q^2)) M_-^2 + (g_0(q^2) - g_\perp(q^2)) q^2]}{q^2 Q_-}, \\ F_T(q^2) &= \tilde{h}_+(q^2), & F_T^V(q^2) &= -\tilde{h}_+(q^2) \frac{M_+}{Q_-} + \tilde{h}_\perp(q^2) \frac{M_-^2 M_+}{q^2 Q_-} - h_\perp(q^2) \frac{M_+}{q^2}, \\ F_T^P(q^2) &= \tilde{h}_+(q^2) \frac{M_-}{Q_-} - \tilde{h}_\perp(q^2) \frac{M_-}{Q_-}, & F_T^S(q^2) &= \tilde{h}_+(q^2) \frac{1}{Q_-} - h_+(q^2) \frac{1}{Q_+} - \tilde{h}_\perp(q^2) \frac{M_-^2}{q^2 Q_+} + h_\perp(q^2) \frac{M_+^2}{q^2 Q_-}. \end{aligned} \quad (21)$$

For the various helicity form factors, we have used the formulas from the latest lattice QCD calculations [97–99], and each form factor can be written as

$$f(q^2) = \frac{1}{1 - q^2/(m_{\text{pole}}^f)^2} [a_0^f + a_1^f z(q^2)], \quad (22)$$

where m_{pole}^f is pole mass, and f represents $f_{+, \perp, 0}$, $g_{+, \perp, 0}$, $h_{+, \perp}$, and $\tilde{h}_{+, \perp}$, respectively. The explicit numerical values of m_{pole}^f , a_0^f , and a_1^f can be found in Refs. [97–99], and we will not list them in the current work. Besides, the definition of the expansion parameter z is

$$z(q^2) = \frac{\sqrt{t_+ - q^2} - \sqrt{t_+ - t_0}}{\sqrt{t_+ - q^2} + \sqrt{t_+ - t_0}}, \quad (23)$$

where $t_+ = (M_{\Lambda_b} + M_{\Lambda_c})^2$ and $t_0 = (M_{\Lambda_b} - M_{\Lambda_c})^2$, respectively.

C. The helicity amplitude

In SM, the semileptonic decay $\Lambda_b \rightarrow \Lambda_c \ell^- \bar{\nu}_\ell$ is considered to be $\Lambda_b \rightarrow \Lambda_c W_{\text{off shell}}^-$, and the off shell $W_{\text{off shell}}^-$ decays to $\ell^- \bar{\nu}_\ell$, subsequently. It is known to us that the $W_{\text{off shell}}^-$ has four helicities, namely $\lambda_W = \pm 1, 0 (J = 1)$ and $\lambda_W = 0 (J = 0)$, and only the off shell $W_{\text{off shell}}^-$ has a timelike polarization, with $J = 1, 0$ denoting the two

FIG. 1. Definition of the polar angles θ , θ_B , and the azimuthal angle χ .

angular momenta of the rest frame $W_{\text{off shell}}^-$. In order to distinguish the two $\lambda_W = 0$ states, we adopt the notation $\lambda_W = 0$ for $J = 1$ and $\lambda_W = t$ for $J = 0$. In the Λ_b -baryon rest frame, choosing the z axis to be along the $W_{\text{off shell}}^-$ (see Fig. 1), we can write the polarization of the $W_{\text{off shell}}^-$ as

$$\begin{aligned} \epsilon^\mu(\pm) &= \frac{1}{\sqrt{2}}(0, 1, \mp i, 0); & \epsilon^\mu(0) &= -\frac{1}{\sqrt{q^2}}(q_z, 0, 0, q_0); \\ \epsilon^\mu(t) &= -\frac{q^\mu}{\sqrt{q^2}}; \end{aligned} \quad (24)$$

where q^μ is the four momentum of the $W_{\text{off shell}}^-$. In this case, the polarization vectors of the $W_{\text{off shell}}^-$ satisfy the orthonormality and completeness relations,

$$\epsilon^{*\mu}(m)\epsilon_\mu(n) = g_{mn}, \quad \sum_{m,n} \epsilon^{*\mu}(m)\epsilon^\nu(n)g_{mn} = g^{\mu\nu} \quad (25)$$

with $g_{mn} = \text{diag}(+, -, -, -)$ for $m, n = t, \pm, 0$. Because the current is not conserved in SM, it consists of a superposition of a spin-1 and a spin-0 component, where the J^P content of the vector current J_μ^V and the axial vector current J_μ^A are $(0^+, 1^-)$ and $(0^-, 1^+)$, respectively. Therefore, we define helicity amplitudes as

$$\begin{aligned} H_{\lambda_2, \lambda_W}^L &= H_{\lambda_2, \lambda_W}^V - H_{\lambda_2, \lambda_W}^A, \\ H_{\lambda_2, \lambda_W}^{V(A)} &= \epsilon^{\dagger\mu}(\lambda_W) \langle \Lambda_c, \lambda_2 | \bar{c} \gamma_\mu (\gamma_5) b | \Lambda_b, \lambda_1 \rangle. \end{aligned} \quad (26)$$

Due to the conservation of angular momentum, $\lambda_1 = -\lambda_2 + \lambda_W$ is satisfied. With above definitions, we preform the helicity amplitudes and have (see, e.g., Refs. [114–116])

$$\begin{aligned} H_{+\frac{1}{2}t}^{V/A} &= \frac{\sqrt{Q_\pm}}{\sqrt{q^2}} \left(M_\mp F_1^{V/A} \pm \frac{q^2}{M_{\Lambda_b}} F_3^{V/A} \right), \\ H_{+\frac{1}{2}+1}^{V/A} &= \sqrt{2Q_\mp} \left(F_1^{V/A} \pm \frac{M_\pm}{M_{\Lambda_b}} F_2^{V/A} \right), \\ H_{+\frac{1}{2}0}^{V/A} &= \frac{\sqrt{Q_\mp}}{\sqrt{q^2}} \left(M_\pm F_1^{V/A} \pm \frac{q^2}{M_{\Lambda_b}} F_2^{V/A} \right). \end{aligned} \quad (27)$$

From parity or from an explicit calculation, we have

$$H_{-\lambda_2, -\lambda_W}^V = H_{\lambda_2, \lambda_W}^V, \quad H_{-\lambda_2, -\lambda_W}^A = -H_{\lambda_2, \lambda_W}^A. \quad (28)$$

If the right-handed current exists, we also have

$$H_{\lambda_2, \lambda_W}^R = H_{\lambda_2, \lambda_W}^V + H_{\lambda_2, \lambda_W}^A. \quad (29)$$

For the $(S \mp P)$ -type current, the corresponding helicity amplitudes are given by [105]

$$H_{\pm\frac{1}{2}0}^{SPL} = \frac{\sqrt{Q_+}}{m_b - m_c} \left(F_1^V M_- + F_3^V \frac{q^2}{M_{\Lambda_b}} \right) \pm \frac{\sqrt{Q_-}}{m_b + m_c} \left(F_1^A M_+ - F_3^A \frac{q^2}{M_1} \right), \quad (30)$$

$$H_{\pm\frac{1}{2}0}^{SPR} = \frac{\sqrt{Q_+}}{m_b - m_c} \left(F_1^V M_- + F_3^V \frac{q^2}{M_{\Lambda_b}} \right) \mp \frac{\sqrt{Q_-}}{m_b + m_c} \left(F_1^A M_+ - F_3^A \frac{q^2}{M_1} \right). \quad (31)$$

Also, one could define the hadronic helicity amplitudes of the (pseudo)-tensor-type current as

$$H_{\lambda_2, \lambda_W, \lambda_{W'}}^T = \epsilon^{\dagger\mu}(\lambda_W) \epsilon^{\dagger\nu}(\lambda_{W'}) \langle \Lambda_c, \lambda_2 | \bar{c} i \sigma_{\mu\nu} (1 - \gamma_5) b | \Lambda_b, \lambda_1 \rangle, \quad (32)$$

and their explicit expressions are given by

$$\begin{aligned}
H_{\frac{1}{2},+,0}^T &= -\sqrt{\frac{2}{q^2}}(F_T\sqrt{Q_+}M_- - F_T^P\sqrt{Q_+}Q_- + G_T\sqrt{Q_-}M_+ + G_T^P\sqrt{Q_-}Q_+), \\
H_{\frac{1}{2},+,-}^T &= -F_T\sqrt{Q_+} - G_T\sqrt{Q_-}, \\
H_{\frac{1}{2},+,t}^T &= -\sqrt{\frac{2}{q^2}}(F_T\sqrt{Q_-}M_+ - F_T^P\sqrt{Q_-}M_+M_- + G_T\sqrt{Q_+}M_- + G_T^P\sqrt{Q_+}M_+M_-) + \sqrt{2q^2}(F_T^V\sqrt{Q_-} - G_T^V\sqrt{Q_+}), \\
H_{\frac{1}{2},0,t}^T &= -F_T\sqrt{Q_-} - G_T\sqrt{Q_+} + F_T^P\sqrt{Q_-}M_- - G_T^P\sqrt{Q_+}M_+ + F_T^V\sqrt{Q_-}M_+ - G_T^V\sqrt{Q_+}M_- + F_T^S\sqrt{Q_-}Q_+ + G_T^S\sqrt{Q_+}Q_-, \\
H_{-\frac{1}{2},+,-}^T &= F_T\sqrt{Q_+} - G_T\sqrt{Q_-}, \\
H_{-\frac{1}{2},0,-}^T &= \sqrt{\frac{2}{q^2}}(F_T\sqrt{Q_+}M_- - F_T^P\sqrt{Q_+}Q_- - G_T\sqrt{Q_-}M_+ - G_T^P\sqrt{Q_-}Q_+), \\
H_{-\frac{1}{2},0,t}^T &= -F_T\sqrt{Q_-} + G_T\sqrt{Q_+} + F_T^P\sqrt{Q_-}M_- + G_T^P\sqrt{Q_+}M_+ + F_T^V\sqrt{Q_-}M_+ + G_T^V\sqrt{Q_+}M_- + F_T^S\sqrt{Q_-}Q_+ - G_T^S\sqrt{Q_+}Q_-, \\
H_{-\frac{1}{2},-,t}^T &= -\sqrt{\frac{2}{q^2}}(F_T\sqrt{Q_-}M_+ - F_T^P\sqrt{Q_-}M_+M_- - G_T\sqrt{Q_+}M_- - G_T^P\sqrt{Q_+}M_+M_-) + \sqrt{2q^2}(F_T^V\sqrt{Q_-} + G_T^V\sqrt{Q_+}). \quad (33)
\end{aligned}$$

D. Observable

With the above hadronic helicity amplitudes, we then write down the twofold differential angular decay distribution as

$$\frac{d^2\Gamma(\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell)}{dq^2 d\cos\theta} = N \left[A_1 + \frac{m_\ell^2}{q^2} (A_2^{VL} + A_2^{VR} + A_2^{LR} + A_2^T) + 2A_3 + \frac{4m_\ell}{\sqrt{q^2}} (A_4^{VS} + A_4^{LT} + A_4^{RT}) + A_5 \right], \quad (34)$$

with

$$N = \frac{G_F^2 |V_{cb}|^2 q^2 |\vec{\mathbf{p}}_2|}{512\pi^3 M_{\Lambda_b}^2} \left(1 - \frac{m_\ell^2}{q^2} \right)^2. \quad (35)$$

In above equation, $|\vec{\mathbf{p}}_2| = \sqrt{Q_+ Q_-}/(2M_{\Lambda_b})$ denotes the Λ_c momentum in the Λ_b rest frame. θ is the polar angle of the lepton, the definition of which can be seen in Fig. 1. The auxiliary functions A_i are given as

$$\begin{aligned}
A_1 &= |1 + C_{VL}|^2 [2\sin^2\theta(|H_{\frac{1}{2},0}^L|^2 + |H_{-\frac{1}{2},0}^L|^2) + (1 - \cos\theta)^2 |H_{\frac{1}{2},+}^L|^2 + (1 + \cos\theta)^2 |H_{-\frac{1}{2},-}^L|^2] \\
&\quad + |C_{VR}|^2 [2\sin^2\theta(|H_{\frac{1}{2},0}^R|^2 + |H_{-\frac{1}{2},0}^R|^2) + (1 - \cos\theta)^2 |H_{\frac{1}{2},+}^R|^2 + (1 + \cos\theta)^2 |H_{-\frac{1}{2},-}^R|^2], \quad (36)
\end{aligned}$$

$$\begin{aligned}
A_2^{VL} &= |1 + C_{VL}|^2 [2\cos^2\theta(|H_{\frac{1}{2},0}^L|^2 + |H_{-\frac{1}{2},0}^L|^2) + \sin^2\theta(|H_{\frac{1}{2},+}^L|^2 + |H_{-\frac{1}{2},-}^L|^2) + 2(|H_{\frac{1}{2},t}^L|^2 + |H_{-\frac{1}{2},t}^L|^2) \\
&\quad - 4\cos\theta(H_{\frac{1}{2},0}^L H_{\frac{1}{2},t}^L + H_{-\frac{1}{2},0}^L H_{-\frac{1}{2},t}^L)], \quad (37)
\end{aligned}$$

$$\begin{aligned}
A_2^{VR} &= |C_{VR}|^2 [2\cos^2\theta(|H_{\frac{1}{2},0}^R|^2 + |H_{-\frac{1}{2},0}^R|^2) + \sin^2\theta(|H_{\frac{1}{2},+}^R|^2 + |H_{-\frac{1}{2},-}^R|^2) + 2(|H_{\frac{1}{2},t}^R|^2 + |H_{-\frac{1}{2},t}^R|^2) \\
&\quad - 4\cos\theta(H_{\frac{1}{2},0}^R H_{\frac{1}{2},t}^R + H_{-\frac{1}{2},0}^R H_{-\frac{1}{2},t}^R)], \quad (38)
\end{aligned}$$

$$\begin{aligned}
A_2^{LR} &= 2\text{Re}[(1 + C_{VL})C_{VR}^*][\sin^2\theta(H_{\frac{1}{2},-}^L H_{-\frac{1}{2},-}^R + H_{\frac{1}{2},+}^L H_{\frac{1}{2},+}^R) + 2\cos^2\theta(H_{-\frac{1}{2},0}^L H_{-\frac{1}{2},0}^R + H_{\frac{1}{2},0}^L H_{\frac{1}{2},0}^R) \\
&\quad - 2\cos\theta(H_{-\frac{1}{2},t}^L H_{-\frac{1}{2},t}^R + H_{-\frac{1}{2},0}^L H_{-\frac{1}{2},t}^R + H_{\frac{1}{2},t}^L H_{\frac{1}{2},0}^R + H_{\frac{1}{2},0}^L H_{\frac{1}{2},t}^R) + 2(H_{-\frac{1}{2},t}^L H_{-\frac{1}{2},t}^R + H_{\frac{1}{2},t}^L H_{\frac{1}{2},t}^R)], \quad (39)
\end{aligned}$$

$$\begin{aligned}
A_2^T &= 4|C_T|^2 [2\sin^2\theta(|H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0,t}^T|^2 + |H_{-\frac{1}{2},+,-}^T + H_{-\frac{1}{2},0,t}^T|^2) \\
&\quad + (1 + \cos\theta)^2 |H_{-\frac{1}{2},0,-}^T + H_{-\frac{1}{2},-,t}^T|^2 + (1 - \cos\theta)^2 |H_{\frac{1}{2},+,0}^T + H_{\frac{1}{2},+,t}^T|^2], \quad (40)
\end{aligned}$$

$$A_3 = 2|C_T|^2[2\cos^2\theta(|H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0,t}^T|^2 + |H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0,t}^T|^2) + \sin^2\theta(|H_{\frac{1}{2},+0}^T + H_{\frac{1}{2},+,t}^T|^2 + |H_{\frac{1}{2},0,-}^T + H_{\frac{1}{2},-,t}^T|^2)] \\ + |C_{SL}|^2(|H_{\frac{1}{2},0}^{SPL}|^2 + |H_{\frac{1}{2},0}^{SPL}|^2) + |C_{SR}|^2(|H_{\frac{1}{2},0}^{SPR}|^2 + |H_{\frac{1}{2},0}^{SPR}|^2), \quad (41)$$

$$A_4^{VS} = \text{Re}[(1 + C_{VL})C_{SL}^*][-\cos\theta(H_{\frac{1}{2},0}^L H_{\frac{1}{2},0}^{SPL} + H_{\frac{1}{2},0}^L H_{\frac{1}{2},0}^{SPL}) + H_{\frac{1}{2},t}^L H_{\frac{1}{2},0}^{SPL} + H_{\frac{1}{2},t}^L H_{\frac{1}{2},0}^{SPL}] \\ + \text{Re}[(1 + C_{VL})C_{SR}^*][-\cos\theta(H_{\frac{1}{2},0}^L H_{\frac{1}{2},0}^{SPR} + H_{\frac{1}{2},0}^L H_{\frac{1}{2},0}^{SPR}) + H_{\frac{1}{2},t}^L H_{\frac{1}{2},0}^{SPR} + H_{\frac{1}{2},t}^L H_{\frac{1}{2},0}^{SPR}] \\ + \text{Re}[C_{VR}C_{SL}^*][-\cos\theta(H_{\frac{1}{2},0}^R H_{\frac{1}{2},0}^{SPL} + H_{\frac{1}{2},0}^R H_{\frac{1}{2},0}^{SPL}) + H_{\frac{1}{2},t}^R H_{\frac{1}{2},0}^{SPL} + H_{\frac{1}{2},t}^R H_{\frac{1}{2},0}^{SPL}] \\ + \text{Re}[C_{VR}C_{SR}^*][-\cos\theta(H_{\frac{1}{2},0}^R H_{\frac{1}{2},0}^{SPR} + H_{\frac{1}{2},0}^R H_{\frac{1}{2},0}^{SPR}) + H_{\frac{1}{2},t}^R H_{\frac{1}{2},0}^{SPR} + H_{\frac{1}{2},t}^R H_{\frac{1}{2},0}^{SPR}] \quad (42)$$

$$A_4^{LT} = -2\text{Re}[(1 + C_{VL})C_T^*]\{(1 + \cos\theta)H_{\frac{1}{2},-}^L(H_{\frac{1}{2},-,t}^T + H_{\frac{1}{2},0,-}^T) + (1 - \cos\theta)(H_{\frac{1}{2},+}^L H_{\frac{1}{2},+,t}^T) \\ + H_{\frac{1}{2},0}^L(H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+,-}^T) + H_{\frac{1}{2},0}^L(H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+,-}^T) + H_{\frac{1}{2},+}^L H_{\frac{1}{2},+,0}^T \\ - \cos\theta[H_{\frac{1}{2},t}^L(H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+,-}^T) + H_{\frac{1}{2},t}^L(H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+,-}^T) + H_{\frac{1}{2},+}^L H_{\frac{1}{2},+,0}^T]\}, \quad (43)$$

$$A_4^{RT} = -2\text{Re}(C_{VR}C_T^*)\{(1 + \cos\theta)H_{\frac{1}{2},-}^R(H_{\frac{1}{2},-,t}^T + H_{\frac{1}{2},0,-}^T) + (1 - \cos\theta)(H_{\frac{1}{2},+}^R H_{\frac{1}{2},+,t}^T) \\ + H_{\frac{1}{2},0}^R(H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+,-}^T) + H_{\frac{1}{2},0}^R(H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+,-}^T) + H_{\frac{1}{2},+}^R H_{\frac{1}{2},+,0}^T \\ - \cos\theta[H_{\frac{1}{2},t}^R(H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+,-}^T) + H_{\frac{1}{2},t}^R(H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+,-}^T) + H_{\frac{1}{2},+}^R H_{\frac{1}{2},+,0}^T]\}, \quad (44)$$

$$A_5 = 8\cos\theta\text{Re}(C_{SL}C_T^*)[H_{\frac{1}{2},0}^{SPL}(H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0,t}^T) + H_{\frac{1}{2},0}^{SPL}(H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0,t}^T)] \\ + 8\cos\theta\text{Re}(C_{SR}C_T^*)[H_{\frac{1}{2},0}^{SPR}(H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0,t}^T) + H_{\frac{1}{2},0}^{SPR}(H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0,t}^T)] \\ + 2\text{Re}[(1 + C_{VL})C_{VR}^*][(1 + \cos\theta)^2 H_{\frac{1}{2},-}^L H_{\frac{1}{2},-}^R + (1 - \cos\theta)^2 H_{\frac{1}{2},+}^L H_{\frac{1}{2},+}^R \\ + 2\sin^2\theta(H_{\frac{1}{2},0}^L H_{\frac{1}{2},0}^R + H_{\frac{1}{2},0}^L H_{\frac{1}{2},0}^R)] + 4\text{Re}(C_{SL}C_{SR}^*)(H_{\frac{1}{2},0}^{SPL} H_{\frac{1}{2},0}^{SPR} + H_{\frac{1}{2},0}^{SPL} H_{\frac{1}{2},0}^{SPR}). \quad (45)$$

When integrating out $\cos\theta$ in Eq. (34), we then obtain the differential decay rate $d\Gamma(\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell)/dq^2$, and the total branching fraction can be given as

$$\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell) = \tau_{\Lambda_b} \int_{m_\ell^2}^{M^2} dq^2 \frac{d\Gamma(\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell)}{dq^2}, \quad (46)$$

τ_{Λ_b} denoting the lifetime of Λ_b baryon. Similar to $\mathcal{R}(D^{(*)})$, one can define the differential and integrated ratios as

$$\mathcal{R}(\Lambda_c)(q^2) = \frac{d\Gamma(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)/dq^2}{d\Gamma(\Lambda_b \rightarrow \Lambda_c \mu \bar{\nu}_\mu)/dq^2}, \quad (47)$$

$$\mathcal{R}(\Lambda_c) = \frac{\int_{m_\tau^2}^{M^2} dq^2 d\Gamma(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)/dq^2}{\int_{m_\mu^2}^{M^2} dq^2 d\Gamma(\Lambda_b \rightarrow \Lambda_c \mu \bar{\nu}_\mu)/dq^2}. \quad (48)$$

Also, the forward-backward asymmetry in the lepton side is defined as

$$A_{\text{FB}}(q^2) = \frac{\int_0^1 d\cos\theta(d^2\Gamma/dq^2 d\cos\theta) - \int_{-1}^0 d\cos\theta(d^2\Gamma/dq^2 d\cos\theta)}{\int_0^1 d\cos\theta(d^2\Gamma/dq^2 d\cos\theta) + \int_{-1}^0 d\cos\theta(d^2\Gamma/dq^2 d\cos\theta)}. \quad (49)$$

Besides, we can calculate the helicity-dependent differential decay rates, which are given as (normalized by the pre-factor $N/2$),

$$\begin{aligned}
\frac{d\Gamma^{\lambda_2=1/2}}{dq^2} &= \frac{m_\ell^2}{q^2} \left[\frac{8}{3} |1 + C_{VL}|^2 (|H_{\frac{1}{2},+}^L|^2 + |H_{\frac{1}{2},0}^L|^2 + 3|H_{\frac{1}{2},t}^L|^2) + \frac{8}{3} |C_{VR}|^2 (|H_{\frac{1}{2},0}^R|^2 + |H_{\frac{1}{2},+}^R|^2 + 3|H_{\frac{1}{2},t}^R|^2) \right. \\
&\quad + \frac{64}{3} |C_T|^2 (|H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0,t}^T|^2 + |H_{\frac{1}{2},+0}^T + H_{\frac{1}{2},+,t}^T|^2) + \frac{16}{3} \text{Re}[(1 + C_{VL})C_{VR}^*] (H_{\frac{1}{2},0}^L H_{\frac{1}{2},0}^R + H_{\frac{1}{2},+}^L H_{\frac{1}{2},+}^R + 3H_{\frac{1}{2},t}^L H_{\frac{1}{2},t}^R) \\
&\quad + \frac{16m_\ell}{\sqrt{q^2}} \text{Re}[(1 + C_{VL})H_{\frac{1}{2},t}^L + C_{VR}H_{\frac{1}{2},t}^R] (C_{SL}^* H_{\frac{1}{2},0}^{SPL} + C_{SR}^* H_{\frac{1}{2},0}^{SPR}) - 2C_T^* [(1 + C_{VL})H_{\frac{1}{2},0}^L + C_{VR}H_{\frac{1}{2},0}^R] (H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+,-}^T) \\
&\quad - 2C_T^* [(1 + C_{VL})H_{\frac{1}{2},+}^L + C_{VR}H_{\frac{1}{2},+}^R] (H_{\frac{1}{2},+0}^T + H_{\frac{1}{2},+,t}^T) \\
&\quad + \frac{16}{3} \left[|1 + C_{VL}|^2 (|H_{\frac{1}{2},0}^L|^2 + |H_{\frac{1}{2},+}^L|^2) + |C_{VR}|^2 (|H_{\frac{1}{2},0}^R|^2 + |H_{\frac{1}{2},+}^R|^2) + \frac{3}{2} |(C_{SL} H_{\frac{1}{2},0}^{SPL} + C_{SR} H_{\frac{1}{2},0}^{SPR})|^2 \right. \\
&\quad \left. + 2|C_T|^2 (|H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+,-}^T|^2 + |H_{\frac{1}{2},+0}^T + H_{\frac{1}{2},+,t}^T|^2) + 2\text{Re}[(1 + C_{VL})C_{VR}^*] (H_{\frac{1}{2},0}^L H_{\frac{1}{2},0}^R + H_{\frac{1}{2},+}^L H_{\frac{1}{2},+}^R) \right], \\
\frac{d\Gamma^{\lambda_2=-1/2}}{dq^2} &= \frac{m_\ell^2}{q^2} \left[\frac{8}{3} |1 + C_{VL}|^2 (|H_{\frac{1}{2},-}^L|^2 + |H_{\frac{1}{2},0}^L|^2 + 3|H_{\frac{1}{2},t}^L|^2) + \frac{8}{3} |C_{VR}|^2 (|H_{\frac{1}{2},-}^R|^2 + |H_{\frac{1}{2},0}^R|^2 + 3|H_{\frac{1}{2},t}^R|^2) \right. \\
&\quad + \frac{64}{3} |C_T|^2 (|H_{\frac{1}{2},+,-}^{T2} + H_{\frac{1}{2},0,t}^{T2}|^2 + |H_{\frac{1}{2},-0}^{T2} + H_{\frac{1}{2},-,t}^{T2}|^2) \\
&\quad + \frac{16}{3} \text{Re}[(1 + C_{VL})C_{VR}^*] (H_{\frac{1}{2},-}^L H_{\frac{1}{2},-}^R + H_{\frac{1}{2},0}^L H_{\frac{1}{2},0}^R + 3H_{\frac{1}{2},t}^L H_{\frac{1}{2},t}^R) \\
&\quad + \frac{16m_\ell}{\sqrt{q^2}} \text{Re}[(1 + C_{VL})^* H_{\frac{1}{2},t}^L + C_{VR}^* H_{\frac{1}{2},t}^R] (C_{SL} H_{\frac{1}{2},0}^{SPL} + C_{SR} H_{\frac{1}{2},0}^{SPR}) \\
&\quad - 2C_T^* [(1 + C_{VL})H_{\frac{1}{2},0}^L + C_{VR}H_{\frac{1}{2},0}^R] (H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0,t}^T) - 2C_T^* [(1 + C_{VL})H_{\frac{1}{2},-}^L + C_{VR}H_{\frac{1}{2},-}^R] (H_{\frac{1}{2},-0}^T + H_{\frac{1}{2},-,t}^T) \\
&\quad + \frac{16}{3} \left[|1 + C_{VL}|^2 (|H_{\frac{1}{2},-}^L|^2 + |H_{\frac{1}{2},0}^L|^2) + |C_{VR}|^2 (|H_{\frac{1}{2},-}^R|^2 + |H_{\frac{1}{2},0}^R|^2) + \frac{3}{2} (|C_{SL}|^2 |H_{\frac{1}{2},0}^{SPL}|^2 + |C_{SR}|^2 |H_{\frac{1}{2},0}^{SPR}|^2) \right. \\
&\quad + 2|C_T|^2 (|H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0,t}^T|^2 + |H_{\frac{1}{2},-0}^T + H_{\frac{1}{2},-,t}^T|^2) + 2\text{Re}[(1 + C_{VL})C_{VR}^*] (H_{\frac{1}{2},-}^L H_{\frac{1}{2},-}^R + H_{\frac{1}{2},0}^L H_{\frac{1}{2},0}^R) \\
&\quad \left. + 3\text{Re}(C_{SL} C_{SR}^*) H_{\frac{1}{2},0}^{SPL} H_{\frac{1}{2},0}^{SPR} \right] \\
\frac{d\Gamma^{\lambda_\tau=1/2}}{dq^2} &= \frac{m_\ell^2}{q^2} \left[|1 + C_{VL}|^2 \left[\frac{8}{3} (|H_{\frac{1}{2},+}^{VL}|^2 + |H_{\frac{1}{2},0}^{VL}|^2 + |H_{\frac{1}{2},-}^{VL}|^2 + |H_{\frac{1}{2},0}^{VL}|^2) + 8(|H_{\frac{1}{2},t}^{VL}|^2 + |H_{\frac{1}{2},t}^{VL}|^2) \right] \right. \\
&\quad + |C_{VR}|^2 \left[\frac{8}{3} (|H_{\frac{1}{2},+}^R|^2 + |H_{\frac{1}{2},0}^R|^2 + |H_{\frac{1}{2},-}^R|^2 + |H_{\frac{1}{2},0}^R|^2) + 8(|H_{\frac{1}{2},t}^R|^2 + |H_{\frac{1}{2},t}^R|^2) \right] \\
&\quad + \frac{16}{3} \text{Re}[(1 + C_{VL})C_{VR}^*] (H_{\frac{1}{2},-}^L H_{\frac{1}{2},-}^R + H_{\frac{1}{2},0}^L H_{\frac{1}{2},0}^R + 3H_{\frac{1}{2},t}^L H_{\frac{1}{2},t}^R + H_{\frac{1}{2},0}^L H_{\frac{1}{2},0}^R + H_{\frac{1}{2},+}^L H_{\frac{1}{2},+}^R + 3H_{\frac{1}{2},t}^L H_{\frac{1}{2},t}^R) \\
&\quad + 8[|C_{SL}|^2 (|H_{\frac{1}{2},0}^{SPL}|^2 + |H_{\frac{1}{2},0}^{SPL}|^2) + |C_{SR}|^2 (|H_{\frac{1}{2},0}^{SPR}|^2 + |H_{\frac{1}{2},0}^{SPR}|^2) + 2\text{Re}(C_{SL} C_{SR}^*) (H_{\frac{1}{2},0}^{SPL} H_{\frac{1}{2},0}^{SPR} + H_{\frac{1}{2},0}^{SPL} H_{\frac{1}{2},0}^{SPR})] \\
&\quad + \frac{32|C_T|^2}{3} (|H_{\frac{1}{2},-,t}^T + H_{\frac{1}{2},0,-}^T|^2 + |H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+,-}^T|^2 + |H_{\frac{1}{2},+0}^T + H_{\frac{1}{2},+,t}^T|^2 + |H_{\frac{1}{2},-0}^T + H_{\frac{1}{2},-,t}^T|^2) \\
&\quad + \frac{16m_\ell}{\sqrt{q^2}} \text{Re}[(1 + C_{VL})H_{\frac{1}{2},t}^L + C_{VR}H_{\frac{1}{2},t}^R] (C_{SL}^* H_{\frac{1}{2},0}^{SPL} + C_{SR}^* H_{\frac{1}{2},0}^{SPR}) \\
&\quad + [(1 + C_{VL})H_{\frac{1}{2},t}^L + C_{VR}H_{\frac{1}{2},t}^R] (C_{SL}^* H_{\frac{1}{2},0}^{SPL} + C_{SR}^* H_{\frac{1}{2},0}^{SPR}) \\
&\quad - \frac{32m_\ell}{3\sqrt{q^2}} [\text{Re}[(1 + C_{VL})C_T^*] (H_{\frac{1}{2},0}^L H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0}^L H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+}^L H_{\frac{1}{2},+0}^T + H_{\frac{1}{2},+}^L H_{\frac{1}{2},+,t}^T + H_{\frac{1}{2},0}^L H_{\frac{1}{2},+,-}^T \\
&\quad + H_{\frac{1}{2},0}^L H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},-}^L H_{\frac{1}{2},-0}^T + H_{\frac{1}{2},-}^L H_{\frac{1}{2},-,t}^T) \\
&\quad + \text{Re}(C_{VR} C_T^*) (H_{\frac{1}{2},0}^R H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0}^R H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+}^R H_{\frac{1}{2},+0}^T + H_{\frac{1}{2},+}^R H_{\frac{1}{2},+,t}^T + H_{\frac{1}{2},0}^R H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0}^R H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},-}^R H_{\frac{1}{2},-0}^T \\
&\quad + H_{\frac{1}{2},-}^R H_{\frac{1}{2},-,t}^T)],
\end{aligned}$$

$$\begin{aligned}
 \frac{d\Gamma^{\lambda_\tau=-1/2}}{dq^2} = & \frac{16}{3} [|1 + C_{VL}|^2 (|H_{\frac{1}{2},+}^{VL}|^2 + |H_{\frac{1}{2},0}^{VL}|^2 + |H_{\frac{1}{2},-}^{VL}|^2 + |H_{-\frac{1}{2},0}^{VL}|^2) + |C_{VR}|^2 (|H_{\frac{1}{2},-}^R|^2 + |H_{-\frac{1}{2},0}^R|^2 + |H_{\frac{1}{2},0}^R|^2 + |H_{\frac{1}{2},+}^R|^2) \\
 & + 2\text{Re}[(1 + C_{VL})C_{VR}^*] (H_{\frac{1}{2},-}^{VL} H_{\frac{1}{2},-}^R + H_{\frac{1}{2},0}^{VL} H_{-\frac{1}{2},0}^R + H_{\frac{1}{2},0}^{VL} H_{\frac{1}{2},0}^R + H_{\frac{1}{2},+}^{VL} H_{\frac{1}{2},+}^R)] \\
 & + \frac{64m_\ell^2}{3q^2} |C_T|^2 (|H_{\frac{1}{2},+,-}^T|^2 + |H_{\frac{1}{2},0,t}^T|^2 + |H_{\frac{1}{2},+0}^T + H_{\frac{1}{2},+t}^T|^2 + |H_{-\frac{1}{2},+,-}^T + H_{-\frac{1}{2},0,t}^T|^2 + |H_{-\frac{1}{2},0,-}^T + H_{-\frac{1}{2},-t}^T|^2) \\
 & - \frac{64m_\ell}{3\sqrt{q^2}} [\text{Re}[(1 + C_{VL})C_T^*] (H_{\frac{1}{2},0}^L H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0}^L H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+}^L H_{\frac{1}{2},+0}^T + H_{\frac{1}{2},+}^L H_{\frac{1}{2},+t}^T + H_{-\frac{1}{2},0}^L H_{-\frac{1}{2},+,-}^T + H_{-\frac{1}{2},0}^L H_{-\frac{1}{2},0,t}^T \\
 & + H_{-\frac{1}{2},-}^L H_{-\frac{1}{2},0,-}^T + H_{-\frac{1}{2},-}^L H_{-\frac{1}{2},-t}^T) + \text{Re}(C_{VR}C_T^*) (H_{\frac{1}{2},0}^R H_{\frac{1}{2},+,-}^T + H_{\frac{1}{2},0}^R H_{\frac{1}{2},0,t}^T + H_{\frac{1}{2},+}^R H_{\frac{1}{2},+0}^T + H_{\frac{1}{2},+}^R H_{\frac{1}{2},+t}^T \\
 & + H_{-\frac{1}{2},0}^R H_{-\frac{1}{2},+,-}^T + H_{-\frac{1}{2},0}^R H_{-\frac{1}{2},0,t}^T + H_{-\frac{1}{2},-}^R H_{-\frac{1}{2},0,-}^T + H_{-\frac{1}{2},-}^R H_{-\frac{1}{2},-t}^T)]. \quad (50)
 \end{aligned}$$

Based on above helicities, we have two observables, the q^2 -dependent longitudinal polarizations of Λ_c baryon and τ lepton, the definitions of which are given, respectively, as

$$P_L^{\Lambda_c}(q^2) = \frac{d\Gamma^{\lambda_2=1/2}/dq^2 - d\Gamma^{\lambda_2=-1/2}/dq^2}{d\Gamma^{\lambda_2=1/2}/dq^2 + d\Gamma^{\lambda_2=-1/2}/dq^2}; \quad (51)$$

$$P_L^\tau(q^2) = \frac{d\Gamma^{\lambda_\tau=1/2}/dq^2 - d\Gamma^{\lambda_\tau=-1/2}/dq^2}{d\Gamma^{\lambda_\tau=1/2}/dq^2 + d\Gamma^{\lambda_\tau=-1/2}/dq^2}. \quad (52)$$

In Eqs. (36)–(45), one finds that there are many quadratic terms of $\cos\theta$. In order to show up these terms, one defines a convexity parameter as

$$C_F^l(q^2) = \frac{1}{d\Gamma/dq^2} \frac{d^2}{d\cos^2\theta} \left(\frac{d^2\Gamma}{dq^2 d\cos\theta} \right). \quad (53)$$

III. RESULTS AND DISCUSSIONS

A. Input parameters

In this section, we shall present the numerical results and discussions only for two cases $\ell^- = \mu^-$ and τ^- . Because the results for the electron are almost the same as those of the muon mode, we will not discuss the electron case any more. In our calculations, the most important nonperturbative parameters are the hadronic transition form factors that have been specified already in the previous section. The other parameters we used are listed here as

$$\begin{aligned}
 M_{\Lambda_b} &= 5.62 \text{ GeV}, & M_{\Lambda_c} &= 2.29 \text{ GeV}, \\
 m_\tau &= 1.78 \text{ GeV}, & m_\mu &= 0.106 \text{ GeV}, \\
 G_F &= 1.66 \times 10^{-5} \text{ GeV}^{-2}, & m_b &= 4.2 \text{ GeV}, \\
 m_c &= 1.28 \text{ GeV}, & \tau_{\Lambda_b} &= 1.466 \text{ ps}. \quad (54)
 \end{aligned}$$

B. SM predictions

Let us present the numerical results of SM firstly. Using the parameters, we calculate the branching fractions of $\Lambda_b \rightarrow \Lambda_c \ell \bar{\nu}_\ell$ as

$$\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \mu \bar{\nu}_\mu)|^{\text{SM}} = (5.34 \pm 0.35) \times 10^{-2}; \quad (55)$$

$$\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)|^{\text{SM}} = (1.78 \pm 0.10) \times 10^{-2}; \quad (56)$$

and the ratio $\mathcal{R}(\Lambda_b)$ is given by

$$\mathcal{R}(\Lambda_b)^{\text{SM}} = 0.33 \pm 0.01. \quad (57)$$

Using the form factors of QCD sum rules [117], the authors in Ref. [100] obtained $\mathcal{R}(\Lambda_b)^{\text{SM}} \in [0.26, 0.30]$, which is a bit smaller than our prediction. In Ref. [94], the authors calculated the $\Lambda_b \rightarrow \Lambda_c$ form factors in the covariant confined quark model and got $\mathcal{R}(\Lambda_b)^{\text{SM}} = 0.29 \pm 0.01$, and this ratio is also smaller than ours. It is obvious that the uncertainty of this ratio is sensitive form factors. That is, the theoretical uncertainty in SM is actually much larger. In order to reduce the uncertainty, more experimental data are needed in future.

In Fig. 2, the q^2 dependence of the differential ratios $d\Gamma/dq^2$, the forward-backward asymmetries of the leptonic side $A_{FB}(q^2)$, the longitudinal polarization components of the Λ_c and leptons [$P_L^h(q^2)$ and $P_L^l(q^2)$], and the convexity parameter $C_F^l(q^2)$ are presented, respectively. As seen from Fig. 2, for the muon mode, the differential ratio shows a steplike behavior when $q^2 = m_\mu^2$. At zero recoil $q^2 = (M_{\Lambda_b} - M_{\Lambda_c})^2$, A_{FB}^ℓ approaches to zero for both μ and τ modes. However, at a large recoil, the A_{FB}^ℓ are -0.12 and -0.35 , for both μ and τ modes, respectively. Furthermore, A_{FB}^μ is positive in most ranges, and the negative value arises from the small mass of muon. Due to the large mass of the τ lepton, the A_{FB}^τ becomes negative when moving away from the zero recoil, and it goes through the zero point at $q^2 = 7.93 \text{ GeV}^2$. For the parameter C_F^ℓ , at the zero recoil limit, both C_F^μ and C_F^τ are zero. At a large recoil range,

FIG. 2. The q^2 dependence of differential ratios dBr/dq^2 , the forward-backward asymmetries of the leptonic side $A_{FB}(q^2)$, the convexity parameters $C_F^l(q^2)$, and the longitudinal polarization components of the Λ_c and leptons in SM. The red (solid) and blue (dashed) lines indicate muon mode and tau mode, respectively.

$C_F^\mu = -1.4$ when $q^2 = 0.4 \text{ GeV}^2$, while C_F^μ changes to zero quickly when $q^2 = m_\mu^2$ due to the effect of the mass. This behavior indicates that the $\cos\theta$ distribution in $q^2 \in [0.4, 11] \text{ GeV}^2$ is strongly parabolic. On the contrary, the C_F^τ is quite small in the whole range, which implies a straight-line behavior of the $\cos\theta$ distribution. For the longitudinal polarization of the Λ_c , it is found that at zero recoil, the longitudinal polarizations are zero for two modes, while they approach to -1 in the large recoil regions. In the last panel, the q^2 dependence of the longitudinal polarization of the charged lepton are displayed. For the muon mode, the mass of muon lepton can be neglected in the Λ_b decays; thus, the muon lepton can be viewed as purely left-handed, and $P_L^\mu \simeq -1$ in whole region. While for the tau mode, the behavior is quite different, $P_L^\tau = -0.52$ at the zero recoil, and $P_L^\tau = 0$ at the maximal recoil point.

C. Model independence

We now present our discussions in the model-independence manner. Motivated by tensions stated in the previous section, many works have been carried out including the effects of NP, by modifying the Wilson coefficients or introducing new operators. As we know, in the quark level, $\Lambda_b \rightarrow \Lambda_c \tau^- \bar{\nu}_\tau$, $B \rightarrow D^{(*)} \tau^- \bar{\nu}_\tau$, $B_c \rightarrow J/\psi(\eta_c) \tau^- \bar{\nu}_\tau$, and $B_c^- \rightarrow \tau^- \bar{\nu}_\tau$ are all induced by the $b \rightarrow c \tau^- \bar{\nu}_\tau$ transition. So, once we assume that $\mathcal{R}(D^{(*)})$ were induced by NP, it should affect the observables of the $\Lambda_b \rightarrow \Lambda_c \tau^- \bar{\nu}_\tau$ decays. In Fig. 3, we at first depict the correlations among the $\mathcal{R}(D)$, $\mathcal{R}(D^*)$, and $\mathcal{R}(\Lambda_c)$ in the presence of each single

NP operator given in Eq. (8). In the calculations, the form factors of $B \rightarrow D^{(*)}$ we employed are from the heavy quark effective theory including higher power and higher order corrections [28]. The horizontal and vertical bands represent the experimental constraints at 1σ confidence level (C.L.) [12]. It is found from the left panel that the $\mathcal{O}_{VL}$, $\mathcal{O}_{VR}$, $\mathcal{O}_{SL}$, and $\mathcal{O}_T$ operators can accommodate the experimental results of $\mathcal{R}(D^{(*)})$ within 1σ , and $\mathcal{O}_{SR}$ will be ruled out. Furthermore, the $\mathcal{O}_{VR}$ covers most regions and has less prediction power, so we will not discuss $\mathcal{O}_{SR}$ and $\mathcal{O}_{VR}$ temporarily. Using these three figures, we can predict the value of $\mathcal{R}(\Lambda_c)$ with the existence of one of three operators. For instance, if the new physics only contributes to the $\mathcal{O}_{VL}$ operator, one can read the $\mathcal{R}(\Lambda_c) \in [0.37, 0.44]$ from the center panel that describes the correlation between $\mathcal{R}(D)$ and $\mathcal{R}(\Lambda_c)$. Also, $\mathcal{R}(\Lambda_c) \in [0.37, 0.41]$ can be arrived based on the relation between $\mathcal{R}(D^*)$ and $\mathcal{R}(\Lambda_c)$. When combining them together, one can predict that $\mathcal{R}(\Lambda_c) \sim [0.37, 0.41]$, which is a bit larger than that of SM. For the operator $\mathcal{O}_{SL}$, the predicted $\mathcal{R}(\Lambda_c)$ is in $[0.31, 0.45]$, and the main constraint is from $\mathcal{R}(D)$. On the contrary, for the operator $\mathcal{O}_T$, the constraint arises mainly from the $\mathcal{R}(D^*)$, and $\mathcal{R}(\Lambda_c)$ is predicted to be in $[0.28, 0.37]$. In this respect, even if we measure $\mathcal{R}(\Lambda_c)$ close to the SM prediction, we cannot exclude the contributions from the $\mathcal{O}_{SL}$ and $\mathcal{O}_T$ operators.

It should be emphasized that in the above discussions, we have not considered the constraints from the recent measurement of $\mathcal{R}(J/\psi)$ and the upper limit of branching fraction of pure leptonic decay $B_c^- \rightarrow \tau^- \bar{\nu}_\tau$. In fact, once these constraints were added, all single operators cannot

FIG. 3. Correlations between $\mathcal{R}(D)$, $\mathcal{R}(D^*)$, and $\mathcal{R}(\Lambda_c)$ in the presence of single NP operators. The vertical and horizontal bands show the experimental constraints, and the black dots denote the SM predictions.

accommodate the 1σ experimental constraints anymore, which have been concluded in Ref. [28]. In order to make a general model-independent analysis, the authors in Ref. [28] had performed a minimum χ^2 fit of the Wilson coefficients to the experimental data of all related observables, namely $\mathcal{R}(D)$, $\mathcal{R}(D^*)$, $\mathcal{R}(J/\psi)$, the longitudinal polarizations of the τ lepton and final meson (D^*), with the single-operator assumption. As aforementioned, the Belle Collaboration released their new results of $\mathcal{R}(D^{(*)})$ [20] at Moriond EW 2019. So, within the latest data, we shall refit the Wilson coefficients including the 2σ uncertainties, and

the updated results are listed in Table I. The technological details are referred to Ref. [28]. It is noted that in contrast to Ref. [28] where all experimental results from *BABAR*, *Belle*, and *LHCb* were input, we here only input the five averaged results. With the new obtained Wilson coefficients, we recalculate $\mathcal{R}(D)$, $\mathcal{R}(D^*)$, and $\mathcal{R}(\Lambda_c)$, and they are also presented in Table I, from which one finds that all operators can enhance the $\mathcal{R}(\Lambda_c)$. Typically, for the operator $\mathcal{O}_{SL}$, the $\mathcal{R}(\Lambda_c)$ can be as large as 0.44.

With the new Wilson coefficients, we also plot the q^2 dependence of the differential branching fraction

TABLE I. Fitted values of the Wilson coefficients with a different NP operator.

NP scenario	Value	χ^2/dof	$\mathcal{R}(D)$	$\mathcal{R}(D^*)$	$\mathcal{R}(\Lambda_c)$
C_{VL}	$(1 + \text{Re}[C_{V_1}])^2 + (\text{Im}[C_{V_1}])^2 = 1.20$	1.51/3	0.34	0.30	0.40
C_{VR}	$-0.002 \pm 0.45i$	1.13/3	0.33	0.30	0.40
C_{SL}	$-0.92 \pm 0.97i$	0.91/3	0.34	0.30	0.44
C_{SR}	0.25	3.40/3	0.39	0.26	0.39
C_T	$-0.005 \pm 0.09i$	2.95/3	0.28	0.30	0.37
SM			0.28	0.25	0.33

$d\mathcal{B}(\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau)/dq^2$ and the ratio $\mathcal{R}(\Lambda_c)$ in Fig. 4. It is seen that these two observables are sensitive to the NP operators. Moreover, the effects of the $\mathcal{O}_{SL}$ operator are remarkable, and one of the other four operators are very similar, although all of them can enhance the branching fraction. Furthermore, we also show in Fig. 5, the q^2 dependence of the Λ_c and τ longitudinal polarizations as

well as the lepton-side forward-backward asymmetry and the convexity parameter $C_F^l(q^2)$. For the $\mathcal{O}_{VL}$ operator, since the resulting Wilson coefficients $(1 + C_{VL})$ appear in both the numerator and the denominator of $P_L^{\Lambda_c, \tau}$ and A_{FB}^τ , the contributions of NP effects are canceled out by each other. As for other operators, due to the interference between SM and NP, the effects of NP can show up.

FIG. 4. Predictions for the differential ratio of $\Lambda_b \rightarrow \Lambda_c \tau \bar{\nu}_\tau$ and $\mathcal{R}(\Lambda_c)$. The solid (black) lines denotes the SM predictions. The dotted (blue), dashed (magenta), dot-dashed (orange), double-dot dashed (brown), and tridot-dashed (red) lines mean NP predictions corresponding to the best-fit Wilson coefficients of $\mathcal{O}_{VL}$, $\mathcal{O}_{VR}$, $\mathcal{O}_{SL}$, $\mathcal{O}_{SR}$, and $\mathcal{O}_T$, respectively.

FIG. 5. Predictions for the forward-backward asymmetry, the convexity parameter $C_F^l(q^2)$, the longitudinal polarization components of the Λ_c , and leptons $P_L^h(q^2)$. The solid (black) lines denote the SM predictions. The dotted (blue), dashed (magenta), dot-dashed (orange), double-dot dashed (brown), and tridot-dashed (red) lines mean NP predictions corresponding to the best-fit Wilson coefficients of $\mathcal{O}_{VL}$, $\mathcal{O}_{VR}$, $\mathcal{O}_{SL}$, $\mathcal{O}_{SR}$, and $\mathcal{O}_T$, respectively.

The effects of $\mathcal{O}_{SL}$ are significant. To be honest, there are certain uncertainties in the hadronic form factors. However, because the uncertainties from the form factor can be largely removed, these four diagrams can be used to be the probes of the new physics, if their observables can be measured precisely in the future, such as in the LHCb or the proposed circular electron-positron collider.

D. Models

In this section, we shall discuss some typical NP models that could affect the $b \rightarrow c\tau\bar{\nu}_\tau$ transition. As we stated previously, the intention of measurements of $\mathcal{R}(D^{(*)})$ is to search for the charged Higgs-like scalars, such as $H^\pm$ in two-Higgs doublet models (2HDMs). As is well known, the charged Higgs boson contributes $b \rightarrow c\tau\bar{\nu}_\tau$ at tree level and its effects are enhanced in some 2HDMs by introducing $\mathcal{O}_{SL}$ and $\mathcal{O}_{SR}$ operators. Unfortunately, these models are disfavored with the combined experimental results; see Ref. [13]. Besides that, models with charged Higgs bosons lead to unexpected large effects in the total B_c lifetime and can also be disfavored by the q^2 distribution. In these respects, we will not discuss the Higg-like models in this current work. Since the charged Higgs effects on $b \rightarrow c\tau\bar{\nu}_\tau$ in the minimal supersymmetric standard model are the same as those in the 2HDM of type II at the tree level, we will not study it, either. As for the R -parity violating minimal supersymmetric standard model, it may affect $b \rightarrow c\tau\bar{\nu}_\tau$ via mediating the slepton and down squark. However, because its behavior is very similar to the leptoquarks that we will discuss later, we will not discuss this model any more. In the following, we mainly focus on the W' model, the scalar leptoquark (LQ) model and the vector LQ model as well.

1. W' model

One simple way to obtain a new physics contribution to $b \rightarrow c\tau\bar{\nu}_\tau$ is to use W' gauge boson which couples to the second- and third-generation fermions [118]. Without the right-handed neutrinos, the general Lagrangian that describes the couplings of an extra W' boson to quarks and leptons can be written as [119]

$$\mathcal{L}_{\text{eff}}^{W'} = \frac{W'}{\sqrt{2}} [\bar{u}_i \gamma^\mu (\epsilon_{u_i d_j}^L P_L + \epsilon_{u_i d_j}^R P_R) d_j + \bar{\ell}_i \epsilon_{\ell_i \nu_j}^L \gamma^\mu P_L \nu_j] + \text{H.c.}, \quad (58)$$

with $u_i \in (u, c, t)$, $d_j \in (d, s, b)$, and $\ell_i, \ell_j \in (e, \mu, \tau)$. $\epsilon_{u_i d_j}^{L,R}$ and $\epsilon_{\ell_i \nu_j}^L$ are the dimensionless coupling parameters that can be constrained by the current data of direct or indirect searches. In addition, for simplicity, all the effective couplings are set to be real. Because we hope the new physics only appears in the τ sector, only the matrix element $\epsilon_{\tau\nu_\tau}^L$ is assumed to be nonzero. If we transfer the

Eq. (58) to the effective Lagrangian as Eq. (7), the Wilson coefficients including the contribution of W' can be expressed as

$$C_{VL} \equiv \frac{\sqrt{2}}{4G_F V_{cb}} \frac{\epsilon_{cb}^L \epsilon_{\tau\nu_\tau}^L}{M_{W'}^2}, \quad C_{VR} \equiv \frac{\sqrt{2}}{4G_F V_{cb}} \frac{\epsilon_{cb}^R \epsilon_{\tau\nu_\tau}^L}{M_{W'}^2}, \quad (59)$$

where $M_{W'}$ is the mass of the W' boson, and $\epsilon_{cb}^{L,R}$ and $\epsilon_{\tau\nu_\tau}^L$ are the flavor-dependent couplings of Eq. (58). As aforementioned, the $\mathcal{O}_{VR}$ cannot contribute to the LFU violation in the first order within SMEFT, so we only discuss the effect of the left-handed W' currently. In fact, when all current experimental constraints are taken into account, the permitted region of $|\epsilon_{cb}^R \epsilon_{\tau\nu_\tau}^L|/M_{W'}^2$ is barely reduced, though the authors of Ref. [120] stated that the right-handed $M_{W'}$ could induce large CP asymmetry in the $B \rightarrow D^* \mu \nu$ decay mode.

To constrain the parameters in Eq. (59), many works have been preformed by analyzing the $\mathcal{R}(D^{(*)})$ anomalies, such as in Refs. [119, 121–123]. In Ref. [121], within the current data of ATLAS [124] and CMS [125], the authors studied $\mathcal{R}(D^{(*)})$ anomalies and the mono-tau signature at the LHC in the left-handed W' model and obtained $\epsilon_{cb}^L \epsilon_{\tau\nu_\tau}^L = (0.14 \pm 0.03)$ with W' mass in the range [0.5, 3.5] TeV, which was consistent with the value $\epsilon_{cb}^L \epsilon_{\tau\nu_\tau}^L = 0.107$ obtained in [122]. Very recently, in Ref. [119], the authors performed a χ^2 analysis within the latest experimental data and got the best fitted values $\epsilon_{cb}^L \epsilon_{\tau\nu_\tau}^L = 0.11$ with $M_{W'} = 1$ TeV, which also agrees with previous studies. So, based on the above discussions and results in Table I, in order to maximize the effects of NP, we adopt the range as

$$\epsilon_{cb}^L \epsilon_{\tau\nu_\tau}^L = (0.12 \pm 0.03) \left(\frac{M_{W'}}{\text{TeV}} \right)^2. \quad (60)$$

With the above parameters, $\mathcal{R}(\Lambda_c)$ is predicted to be 0.38 ± 0.03 . Furthermore, we show $\mathcal{R}(\Lambda_c)$ as a function of $\epsilon_{cb}^L \epsilon_{\tau\nu_\tau}^L$ in Fig. 6, where the q^2 dependence of the differential branching fraction is also presented. It is found that the curves with an extra W' boson are above ones from SM. Since here we only introduce a left-handed W' boson, it only affects the Wilson coefficient C_{VL} and cannot produce other operators, so that other observables, such as A_{FB} , $P_L^{h,l}$, and C_F^l , are unchanged, compared with the results of SM. In addition, because here we suppose W' only couples to the third family fermions, it cannot be produced in the current proton-antiproton collider, so it can escape from the current constrain of direct searches from ATLAS and CMS at LHC.

FIG. 6. The prediction of $\mathcal{R}(\Lambda_c)$ as a function of $\epsilon_{cb}^L \epsilon_{\tau\nu}^L$ (left panel), and q^2 variation of the differential branching fraction (right panel) in the left-handed W' model. The black and red lines indicate the results of SM and the left-handed W' model, respectively. The yellow band is from the uncertainties of W' model.

2. Scalar leptoquark model

For the $b \rightarrow c \ell \bar{\nu}_\ell$, another potential tree-level mediators are leptoquarks. In the literature, there are about ten possible leptoquark models that respect the SM symmetry, assuming the leptoquark be a scalar, a vector, or a tensor particle. For the scalar particle, we discuss here a model [58], where the LQ ϕ transforms under the SM gauge group as $(\mathbf{3}, \mathbf{1}, -\frac{1}{3})$. It is in this model that not only the scalar and vector operators but the tensor operator can be generated. The part of the Lagrangian relevant to the couplings between LQ and SM fermions is given by [58]

$$\mathcal{L}_{\text{int}}^\phi \supset \bar{Q}_L^c \lambda^L i \tau_2 L \phi^* + \bar{u}_R^c \lambda^R \ell_R \phi^* + \text{H.c.}, \quad (61)$$

where Q_L , L denote the left-handed quark and lepton doublet, and u_R , ℓ_R denote the right-handed up-type quark and lepton singlet, respectively. $\lambda^{L,R}$ are the Yukawa coupling matrices in flavor space. More detailed discussions can be found in Ref. [58]. Such a scalar LQ ϕ could induce the $b \rightarrow c \tau \bar{\nu}_\tau$ transition with operators $\mathcal{O}_{VL}$, $\mathcal{O}_{SL}$, and $\mathcal{O}_T$, and the corresponding Wilson coefficients at a scale $\mu = M_\phi$ are given by

$$\begin{aligned} C_{VL}(M_\phi) &= \frac{\lambda_{b\nu_\tau}^L \lambda_{c\tau}^{L*}}{4\sqrt{2}G_F V_{cb} M_\phi^2}, \\ C_{SL}(M_\phi) &= -\frac{\lambda_{b\nu_\tau}^L \lambda_{c\tau}^{R*}}{4\sqrt{2}G_F V_{cb} M_\phi^2}, \\ C_T(M_\phi) &= -\frac{1}{4} C_{SL}(M_\phi). \end{aligned} \quad (62)$$

Because the decay we are concerned with is at a scale $\mu = m_b$, all Wilson coefficients should be run down to $\mu = m_b$ scale. Due to the conservation of the vector current, C_{VL} is not renormalized. In the cases of C_{SL} and C_T , the explicit evolution equations and discussions could be found in Refs. [75,126]. It should be emphasized that the model

with such a scalar LQ can also accommodate the $\mathcal{R}_K$ and $(g-2)_\mu$ anomalies. Taking $M_\phi = 1$ TeV as a benchmark, in Refs. [57,58,61,75], the authors had fitted the parameters within the experimental data; one point they favored is $\lambda_{b\nu_\tau}^L \lambda_{c\tau}^{L*} = 0.35$ and $\lambda_{b\nu_\tau}^L \lambda_{c\tau}^{R*} = -0.03$. Therefore, to maximize the effect of LQ, considering the fitted results in Table. I, we set the parameter space as

$$\begin{aligned} \lambda_{b\nu_\tau}^L \lambda_{c\tau}^{L*} &= (0.35 \pm 0.05) \left(\frac{M_\phi}{\text{TeV}} \right)^2, \\ \lambda_{b\nu_\tau}^L \lambda_{c\tau}^{R*} &= (-0.03 \pm 0.03) \left(\frac{M_\phi}{\text{TeV}} \right)^2. \end{aligned} \quad (63)$$

Using above parameters, we get $\mathcal{R}(\Lambda_c) = 0.41 \pm 0.02$, which is larger than the SM prediction by 25%. By setting $M_\phi = 1$ TeV, we plot the changes of $\mathcal{R}(\Lambda_c)$ with $\lambda_{b\nu_\tau}^L \lambda_{c\tau}^{L*}$ in Fig. 7, and the differential branching fraction as a function of q^2 is also displayed.

3. Vector leptoquark model

Now, we move to discuss the effect of the vector leptoquark. In this article, we shall explore a model [59] where the introduced vector $SU(2)_L$ triplet LQ U_3^μ transforms as $(\mathbf{3}, \mathbf{3}, \frac{2}{3})$ under the SM gauge group. In the mass basis, the part of Lagrangian is given by

$$\begin{aligned} \mathcal{L}_{U_3} &\supset U_{3\mu}^{(2/3)} [(Vg)_{ij} \bar{u}_i \gamma^\mu P_L \nu_j - g_{ij} \bar{d}_i \gamma^\mu P_L \ell_j] \\ &\quad + U_{3\mu}^{(5/3)} (\sqrt{2} Vg)_{ij} \bar{u}_i \gamma^\mu P_L \ell_j \\ &\quad + U_{3\mu}^{(-1/3)} (\sqrt{2} g)_{ij} \bar{d}_i \gamma^\mu P_L \nu_j + \text{H.c.}, \end{aligned} \quad (64)$$

where $i, j = 1, 2, 3$ denote the i, j th generation fermion. The parameter g_{ij} are the couplings of the $Q = 2/3$ component of the triplet, $U_\mu^{(2/3)}$, to $\bar{d}_{Li}$ and ℓ_{Lj} , and V represents the CKM matrix. From the above Lagrangian, it

FIG. 7. The prediction of $\mathcal{R}(\Lambda_c)$ as a function of $\lambda_{b\tau}^L \lambda_{c\tau}^{L*}$ (left panel) and the differential branching fraction as a function of q^2 (right panel) with a scalar leptoquark. The black and red lines indicate the results of SM and the scalar leptoquark model, respectively. The yellow band is from the uncertainties of NP model.

FIG. 8. The prediction of $\mathcal{R}(\Lambda_c)$ as a function of $\lambda_{b\tau}^L \lambda_{c\tau}^{L*}$ (left panel) and the differential branching fraction as a function of q^2 (right panel) with a vector leptoquark. The black and red lines indicate the results of SM and the vector leptoquark model, respectively. The yellow band is from the uncertainties of NP model.

is found that the $b \rightarrow c\tau\bar{\nu}_\tau$ transition proceeds via exchange of LQ with the operator $\mathcal{O}_{VL}$. The corresponding Wilson coefficient C_{VL} is given correspondingly by

$$C_{VL} = \frac{\sqrt{2}g_{b\tau}^*(Vg)_{c\tau}}{4G_F V_{cb} M_U^2}. \quad (65)$$

It should be noted that in this model, due to the existence of the second term the $\mathcal{R}(D^{(*)})$, $\mathcal{R}_{K^{(*)}}$ as well as the angular observable $\mathcal{P}'_5$ in $B \rightarrow K^*\mu^+\mu^-$ decay can be explained simultaneously with suitable parameter space, as shown in Ref. [59]. From the results of χ^2 fits to the measured ratios $\mathcal{R}(D^{(*)})$ and acceptable q^2 spectra done in Ref. [61], we learn that at 1σ we have the following two best-fit solution (S1 and S2):

$$g_{b\tau}^*(Vg)_{c\tau} = \left(\frac{M_U}{\text{TeV}}\right)^2 \begin{cases} 0.18 \pm 0.04, & S_1; \\ -2.88 \pm 0.04, & S_2, \end{cases} \quad (66)$$

which is also in agreement with our fitted results. It can be seen that the vector LQ only affects the functions A_1 and

A_2^{VL} in Eq. (34), and both two terms are proportional to $|1 + C_{VL}|^2$. Therefore, although the fit results S_1 and S_2 are quite different, the coefficients $(1 + C_{VL})$ they induced have nearly same absolute values with different signs. Currently, it is very hard for us to differentiate these two solutions. With the above results, we obtain $\mathcal{R}(\Lambda_c) = 0.43 \pm 0.03$. Also, we plot changes of the $\mathcal{R}(\Lambda_c)$ with $\lambda_{b\tau}^L \lambda_{c\tau}^{L*}$, as well as the differential branching fraction as a function of q^2 in Fig. 8. From the two figures, we find that both the $\mathcal{R}(\Lambda_c)$ and the branching fractions are enhanced by about 30%. From the above analyses, one finds that although the left-handed W' and the vector leptoquark particle origin form different models, however for the concerned decay mode, their behaviors are very similar, since both of them only contribute to the operator $\mathcal{O}_{VL}$.

IV. SUMMARY

Recent measurements of the ratios $\mathcal{R}(D^{(*)})$ imply that a new physics may appear in the $b \rightarrow c\tau\bar{\nu}_\tau$ transition. So, in this work, we explored the effect of NP in the $\Lambda_b \rightarrow \Lambda_c\tau\bar{\nu}_\tau$

decay mode that is also induced by $b \rightarrow c\tau\bar{\nu}_\tau$. To show up the NP effect, we firstly calculated the branching fraction of $\Lambda_b \rightarrow \Lambda_c\tau\bar{\nu}_\tau$ and its ratio to $\Lambda_b \rightarrow \Lambda_c\mu\bar{\nu}_\mu$ in SM, and $\mathcal{R}(\Lambda_c) = 0.33 \pm 0.01$, within the latest form factors from lattice QCD calculations. When studying the effect of NP, we explore $\Lambda_b \rightarrow \Lambda_c\tau\bar{\nu}_\tau$ in a model-independent manner. The differential and total branching fraction, the longitudinal polarizations of final lepton, the forward-backward asymmetries in the lepton side, the convexity parameters, and the ratio $\mathcal{R}(\Lambda_c)$ have been calculated. Supposing that NP only affects the third generation fermion, we also presented the correlations among $\mathcal{R}(D)$, $\mathcal{R}(D^*)$, and $\mathcal{R}(\Lambda_c)$. Considering the latest experimental data, we perform a minimum χ^2 fit of the Wilson coefficient of each operator and found that the left-handed scalar operator $\mathcal{O}_{SL}$ could enhance the branching fraction by 30%. Using the fitted values of the Wilson coefficients, we also give a prognosis for the physical observables of $\Lambda_b \rightarrow \Lambda_c\tau\bar{\nu}_\tau$, such as the ratio $\mathcal{R}(\Lambda_c)$, forward-backward

asymmetry, and other polarized observables as well as the differential branching fraction. At last, we also study the contributions of three typical NP models, the W' model, the scalar leptoquark (LQ) model, and the vector LQ model, to the ratio and the differential branching fraction of $\Lambda_b \rightarrow \Lambda_c\tau\bar{\nu}_\tau$. These results can be tested in the current LHCb experiment and the future high energy experiments.

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