Phenomenology of $\Xi_b \rightarrow \Xi_c \tau \nu$ decays

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Deviations from the standard model prediction have been reported in various semileptonic B decays mediated via $b \rightarrow c$ charged-current interactions. In this context, we analyze corresponding semileptonic baryon decays $\Xi_b \rightarrow \Xi_c \tau \nu$ using the helicity formalism. We report numerical results on various observables such as the decay rate, ratio of branching ratios, lepton-side forward-backward asymmetry, longitudinal polarization fraction of the charged lepton, and the convexity parameter for this decay mode using results from the relativistic quark model. We also provide an estimate of the new physics effect on these observables under various new physics scenarios.

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I. INTRODUCTION

Lepton flavor universality (LFU) violation has been reported in various semileptonic B meson decays mediated via $b \rightarrow c$ charged-current and $b \rightarrow sl^+l^-$ neutral-current interactions. A combined excess of about 4.1σ from the standard model (SM) prediction has been reported by the Heavy Flavor Averaging Group [1] for R_D and R_{D^*} , where $R_{D^{(*)}}$ represents the ratio of branching ratios of $B \rightarrow D^{(*)} \tau \nu$ to the corresponding $B \rightarrow D^{(*)} l \nu$ decays. Similarly, significant deviation from the SM expectation is observed in the LFU ratios $R_K^{(*)} = \mathcal{B}(B \rightarrow K^{(*)} \mu \mu) / \mathcal{B}(B \rightarrow K^{(*)} e e)$ mediated via flavor-changing neutral-current decays. Measurement of R_K in the range $1 < q^2 < 6 \text{ GeV}^2$ deviates from the SM prediction at the 2.6σ level [2]. Similarly, the measured value of R_{K^*} in the dilepton invariant mass $q^2 = [0.045, 1.1] \text{ GeV}^2$ and $[1.1, 6.0] \text{ GeV}^2$ [3] deviates from the SM prediction at approximately 2.1σ and 2.4σ , respectively. Very recently, the LHCb Collaboration [4] has measured the value of the ratio of the branching ratios $R_{J/\Psi} = \mathcal{B}(B_c \rightarrow J/\Psi \tau \nu) / \mathcal{B}(B_c \rightarrow J/\Psi \mu \nu)$ to be $0.71 \pm 0.17 \pm 0.18$. Comparing this measured value with the SM prediction [5–7], we find the discrepancy to be more than 2σ .

Inspired by the anomalies present in the meson decays mediated via $b \rightarrow c$ charged-current interactions, we study the corresponding baryon decays $\Xi_b \rightarrow \Xi_c \tau \nu$

within the SM and within various new physics (NP) scenarios using the $\Xi_b \rightarrow \Xi_c$ form factors obtained from the relativistic quark model. The $\Xi_b \rightarrow \Xi_c e \nu$ decays have been studied by various authors [8–15]. In this paper, we adopt a model-independent approach, in which we use a general effective Lagrangian for the $b \rightarrow c l \nu$ transition in the presence of NP to investigate the impact of various NP operators on different physical observables, such as the decay rate, ratio of the branching ratios, lepton-side forward-backward asymmetry, longitudinal polarization fraction of the charged lepton, and the convexity parameter for the $\Xi_b \rightarrow \Xi_c \tau \nu$ decays. We first constrain the NP operators using 1σ constraints coming from the measured values of the ratio of branching ratios R_D and R_{D^*} , and then analyze their effects on various observables. Moreover, it was reported in Ref. [16] that the lifetime of the B_c meson puts a severe constraint on scalar-type NP interactions. Calculations of the mean lifetime of the B_c meson $\tau_{B_c} = 0.52_{-0.12}^{+0.18} \text{ ps}$ in the SM [17–19] suggest that $\mathcal{B}(B_c \rightarrow \tau \nu)$ cannot exceed 5% to comply with the current world average of the B_c lifetime. Although this constraint can be relaxed up to about 30%, very recently [20] a more stringent limit of $\mathcal{B}(B_c \rightarrow \tau \nu) \leq 10\%$ was obtained by analyzing the LEP data at the Z peak. The constraint coming from the B_c meson decay width is also discussed in detail. However, we do not use the constraint coming from the $R_{J/\Psi}$ measurement as the error associated with it is rather large. It should be noted that, although the analysis for the $\Xi_b \rightarrow \Xi_c l \nu$ decays presented here is quite similar to the one for the semileptonic $\Lambda_b \rightarrow (\Lambda_c, p) l \nu$ decays of Ref. [21], we do not assume any correlation of NP effects in $b \rightarrow u \tau \nu$ and $b \rightarrow c \tau \nu$ transitions and hence do not include the constraint coming from R_π^l in our present analysis.

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Our paper is organized as follows. In Sec. II, we start with the most general effective Lagrangian in the presence of NP for the $b \rightarrow c\tau\nu$ decays that is valid at the renormalization scale $\mu = m_b$. A brief discussion on $\Xi_b \rightarrow \Xi_c$ transition form factors is also presented. In Sec. II, we write down the helicity amplitudes and define several observables, such as the decay rate, ratio of branching ratios, τ polarization fraction, forward-backward asymmetry, and the convexity parameter for the $\Xi_b \rightarrow \Xi_c l\nu$ decays. In Sec. III, we present our numerical results for all the observables defined in Sec. II. Finally, we present a brief summary of our results and conclude in Sec. IV.

II. EFFECTIVE LAGRANGIAN, HEAVY BARYON FORM FACTORS, AND HELICITY AMPLITUDES

A. Effective weak Lagrangian

The natural way to introduce NP effects in a model-independent approach is to construct an effective Lagrangian for the weak decays that includes both SM and beyond-the-SM physics. We pattern our analysis after that of Ref. [21], and adopt a common notation. In particular, we employ the most general effective weak Lagrangian for the $b \rightarrow cl\nu$ transition decays in the presence of NP valid at the renormalization scale $\mu = m_b$ for the computation of the $\Xi_b \rightarrow \Xi_c l\nu$ decay amplitude [22,23]. That is,

$$\begin{aligned} \mathcal{L}_{\text{eff}} = & -\frac{4G_F}{\sqrt{2}} V_{cb} \{ (1 + V_L) \bar{l}_L \gamma_\mu \nu_L \bar{c}_L \gamma^\mu b_L + V_R \bar{l}_L \gamma_\mu \nu_L \bar{c}_R \gamma^\mu b_R + \tilde{V}_L \bar{l}_R \gamma_\mu \nu_R \bar{c}_L \gamma^\mu b_L + \tilde{V}_R \bar{l}_R \gamma_\mu \nu_R \bar{c}_R \gamma^\mu b_R \\ & + S_L \bar{l}_R \nu_L \bar{c}_R b_L + S_R \bar{l}_R \nu_L \bar{c}_L b_R + \tilde{S}_L \bar{l}_L \nu_R \bar{c}_R b_L + \tilde{S}_R \bar{l}_L \nu_R \bar{c}_L b_R + T_L \bar{l}_R \sigma_{\mu\nu} \nu_L \bar{c}_R \sigma^{\mu\nu} b_L \\ & + \tilde{T}_L \bar{l}_L \sigma_{\mu\nu} \nu_R \bar{c}_L \sigma^{\mu\nu} b_R \} + \text{H.c.}, \end{aligned} \quad (1)$$

where G_F is the Fermi constant, V_{cb} is the relevant Cabibbo-Kobayashi-Maskawa (CKM) matrix element, and $(c, b, l, \nu)_{R,L} = (\frac{1 \pm \gamma_5}{2})(c, b, l, \nu)$. We note that the NP couplings $V_{L,R}$, $S_{L,R}$, and T_L characterizing the NP contributions coming from new vector, scalar, and tensor interactions are associated with left-handed neutrinos, whereas the NP couplings $\tilde{V}_{L,R}$, $\tilde{S}_{L,R}$, and $\tilde{T}_L$ are associated with right-handed neutrinos. In this study, we assume that NP contributions come from new vector and scalar interactions only. We neglect the contributions coming from NP couplings that involve right-handed neutrinos, i.e., $\tilde{V}_{L,R} = \tilde{S}_{L,R} = \tilde{T}_L = 0$. We also assume that NP effects appear in the tau mode only. All of the NP couplings are assumed to be real for our analysis. With these assumptions, we obtain

$$\begin{aligned} \mathcal{L}_{\text{eff}} = & -\frac{G_F}{\sqrt{2}} V_{cb} \{ G_V \bar{l} \gamma_\mu (1 - \gamma_5) \nu_l \bar{c} \gamma^\mu b - G_A \bar{l} \gamma_\mu (1 - \gamma_5) \nu_l \bar{c} \gamma^\mu \gamma_5 b + G_S \bar{l} (1 - \gamma_5) \nu_l \bar{c} b - G_P \bar{l} (1 - \gamma_5) \nu_l \bar{c} \gamma_5 b \} \\ & + \text{H.c.}, \end{aligned} \quad (2)$$

where

$$G_V = 1 + V_L + V_R, \quad G_A = 1 + V_L - V_R, \quad G_S = S_L + S_R, \quad G_P = S_L - S_R.$$

In the SM, one has $V_{L,R} = S_{L,R} = 0$.

B. $\Xi_b \rightarrow \Xi_c$ transition form factors

The hadronic matrix elements of vector and axial-vector currents between two spin- $\frac{1}{2}$ baryons are parametrized in terms of various hadronic form factors as follows:

$$\begin{aligned} M_\mu^V &= \langle B_2, \lambda_2 | J_\mu^V | B_1, \lambda_1 \rangle = \bar{u}_2(p_2, \lambda_2) [f_1(q^2) \gamma_\mu + i f_2(q^2) \sigma_{\mu\nu} q^\nu + f_3(q^2) q_\mu] u_1(p_1, \lambda_1), \\ M_\mu^A &= \langle B_2, \lambda_2 | J_\mu^A | B_1, \lambda_1 \rangle = \bar{u}_2(p_2, \lambda_2) [g_1(q^2) \gamma_\mu + i g_2(q^2) \sigma_{\mu\nu} q^\nu + g_3(q^2) q_\mu] \gamma_5 u_1(p_1, \lambda_1), \end{aligned} \quad (3)$$

where $q^\mu = (p_1 - p_2)^\mu$ is the four-momentum transfer, λ_1 and λ_2 are the helicities of the parent and daughter baryons, respectively, and $\sigma_{\mu\nu} = \frac{i}{2} [\gamma_\mu, \gamma_\nu]$. Here, B_1 represents the Ξ_b baryon and B_2 represents the Ξ_c baryon. When both baryons are heavy, it is also convenient to parametrize the matrix element in terms of the four-velocities v^μ and v'^μ as follows:

$$\begin{aligned} M_\mu^V &= \langle B_2, \lambda_2 | J_\mu^V | B_1, \lambda_1 \rangle = \bar{u}_2(p_2, \lambda_2) [F_1(\omega) \gamma_\mu + F_2(\omega) v_\mu + F_3(\omega) v'_\mu] u_1(p_1, \lambda_1), \\ M_\mu^A &= \langle B_2, \lambda_2 | J_\mu^A | B_1, \lambda_1 \rangle = \bar{u}_2(p_2, \lambda_2) [G_1(\omega) \gamma_\mu + G_2(\omega) v_\mu + G_3(\omega) v'_\mu] \gamma_5 u_1(p_1, \lambda_1), \end{aligned} \quad (4)$$

where $\omega = v \cdot v' = (m_{B_1}^2 + m_{B_2}^2 - q^2)/(2m_{B_1}m_{B_2})$ and m_{B_1} and m_{B_2} are the masses of the B_1 and B_2 baryons, respectively. The two sets of form factors are related via

$$\begin{aligned}
f_1(q^2) &= F_1(q^2) + (m_{B_2} + m_{B_1}) \left[\frac{F_2(q^2)}{2m_{B_1}} + \frac{F_3(q^2)}{2m_{B_2}} \right], \\
f_2(q^2) &= \frac{F_2(q^2)}{2m_{B_1}} + \frac{F_3(q^2)}{2m_{B_2}}, \\
f_3(q^2) &= \frac{F_2(q^2)}{2m_{B_1}} - \frac{F_3(q^2)}{2m_{B_2}}, \\
g_1(q^2) &= G_1(q^2) - (m_{B_2} - m_{B_1}) \left[\frac{G_2(q^2)}{2m_{B_1}} + \frac{G_3(q^2)}{2m_{B_2}} \right], \\
g_2(q^2) &= \frac{G_2(q^2)}{2m_{B_1}} + \frac{G_3(q^2)}{2m_{B_2}}, \\
g_3(q^2) &= \frac{G_2(q^2)}{2m_{B_1}} - \frac{G_3(q^2)}{2m_{B_2}}. \tag{5}
\end{aligned}$$

We use the equation of motion to find the hadronic matrix elements of scalar and pseudoscalar currents. That is,

$$\begin{aligned}
\langle B_2, \lambda_2 | \bar{c}b | B_1, \lambda_1 \rangle &= \bar{u}_2(p_2, \lambda_2) \left[f_1(q^2) \frac{\not{q}}{m_b - m_c} \right. \\
&\quad \left. + f_3(q^2) \frac{q^2}{m_b - m_c} \right] u_1(p_1, \lambda_1), \\
\langle B_2, \lambda_2 | \bar{c}\gamma_5 b | B_1, \lambda_1 \rangle &= \bar{u}_2(p_2, \lambda_2) \left[-g_1(q^2) \frac{\not{q}}{m_b + m_c} \right. \\
&\quad \left. - g_3(q^2) \frac{q^2}{m_b + m_c} \right] \gamma_5 u_1(p_1, \lambda_1), \tag{6}
\end{aligned}$$

where m_b is the mass of the b quark and m_c is the mass of the c quark evaluated at the renormalization scale $\mu = m_b$. For the various invariant form factors F_i and G_i , we follow Ref. [8]. The relevant equations for our calculation are as follows:

$$\begin{aligned}
F_1(\omega) &= \zeta(\omega) + \left(\frac{\bar{\Lambda}}{2m_b} + \frac{\bar{\Lambda}}{2m_c} \right) [2\chi(\omega) + \zeta(\omega)], \\
G_1(\omega) &= \zeta(\omega) + \left(\frac{\bar{\Lambda}}{2m_b} + \frac{\bar{\Lambda}}{2m_c} \right) \left[2\chi(\omega) + \frac{\omega - 1}{\omega + 1} \zeta(\omega) \right], \\
F_2(\omega) &= G_2(\omega) = -\frac{\bar{\Lambda}}{2m_c} \frac{2}{\omega + 1} \zeta(\omega), \\
F_3(\omega) &= -G_3(\omega) = -\frac{\bar{\Lambda}}{2m_b} \frac{2}{\omega + 1} \zeta(\omega), \tag{7}
\end{aligned}$$

where $\bar{\Lambda} = m_{B_1} - m_b$ is the difference between the baryon and heavy-quark masses in the heavy-quark limit $m_b \rightarrow \infty$. Here $\zeta(\omega)$ denotes the Isgur-Wise function. The additional function $\chi(\omega)$ appears due to the $1/m_b$ correction to the heavy-quark effective theory Lagrangian. Near the zero recoil point of the final baryon $\omega = 1$, the functions $\zeta(\omega)$ and $\chi(\omega)$ can be expressed as

$$\begin{aligned}
\zeta(\omega) &= 1 - \rho_\zeta^2(\omega - 1) + c_\zeta(\omega - 1)^2, \\
\chi(\omega) &= \rho_\chi^2(\omega - 1) + c_\chi(\omega - 1)^2, \tag{8}
\end{aligned}$$

where ρ_ζ^2 and c_ζ represent the slope and the curvature of the Isgur-Wise functions, respectively. We refer to Ref. [8] for all the omitted details.

C. Helicity amplitudes

We now proceed to discuss the helicity amplitudes for the baryonic $b \rightarrow c \ell \nu$ decay mode. The helicity amplitudes can be defined by [24–26]

$$H_{\lambda_2 \lambda_W}^{V/A} = M_\mu^{V/A}(\lambda_2) \epsilon^{\dagger\mu}(\lambda_W), \tag{9}$$

where λ_2 and λ_W denote the helicities of the daughter baryon and $W_{\text{off-shell}}^-$, respectively. There are total four helicities $\lambda_W = \pm 1, 0$ and $\lambda_W = 0$ corresponding to the angular momenta $J = 1$ and $J = 0$ of the rest frame $W_{\text{off-shell}}^-$. The convention is to use $\lambda_W = t$ for $J = 0$ in order to distinguish it from $\lambda_W = 0$ for the $J = 1$ state. In the rest frame of the parent baryon B_1 , the helicity amplitudes can be written as [21,27,28]

$$\begin{aligned}
H_{\frac{1}{2}0}^V &= G_V \frac{\sqrt{Q_-}}{\sqrt{q^2}} [(m_{B_1} + m_{B_2})f_1(q^2) - q^2 f_2(q^2)], \\
H_{\frac{1}{2}0}^A &= G_A \frac{\sqrt{Q_+}}{\sqrt{q^2}} [(m_{B_1} - m_{B_2})g_1(q^2) + q^2 g_2(q^2)], \\
H_{\frac{1}{2}1}^V &= G_V \sqrt{2Q_-} [-f_1(q^2) + (m_{B_1} + m_{B_2})f_2(q^2)], \\
H_{\frac{1}{2}1}^A &= G_A \sqrt{2Q_+} [-g_1(q^2) - (m_{B_1} - m_{B_2})g_2(q^2)], \\
H_{\frac{1}{2}t}^V &= G_V \frac{\sqrt{Q_+}}{\sqrt{q^2}} [(m_{B_1} - m_{B_2})f_1(q^2) + q^2 f_3(q^2)], \\
H_{\frac{1}{2}t}^A &= G_A \frac{\sqrt{Q_-}}{\sqrt{q^2}} [(m_{B_1} + m_{B_2})g_1(q^2) - q^2 g_3(q^2)], \tag{10}
\end{aligned}$$

where $Q_\pm = (m_{B_1} \pm m_{B_2})^2 - q^2$. Either from parity or from explicit calculation, one can show that $H_{-\lambda_2 - \lambda_W}^V = H_{\lambda_2 \lambda_W}^V$ and $H_{-\lambda_2 - \lambda_W}^A = -H_{\lambda_2 \lambda_W}^A$. Similarly, the scalar and pseudoscalar helicity amplitudes associated with the NP couplings G_S and G_P can be written as [21,27,28]

$$\begin{aligned}
H_{\frac{1}{2}0}^{SP} &= H_{\frac{1}{2}0}^S - H_{\frac{1}{2}0}^P, \\
H_{\frac{1}{2}0}^S &= G_S \frac{\sqrt{Q_+}}{m_b - m_{q'}} [(m_{B_1} - m_{B_2})f_1(q^2) + q^2 f_3(q^2)], \\
H_{\frac{1}{2}0}^P &= G_P \frac{\sqrt{Q_-}}{m_b + m_{q'}} [(m_{B_1} + m_{B_2})g_1(q^2) - q^2 g_3(q^2)]. \tag{11}
\end{aligned}$$

Moreover, we have $H_{\lambda_2\lambda_{\text{NP}}}^S = H_{-\lambda_2-\lambda_{\text{NP}}}^S$ and $H_{\lambda_2\lambda_{\text{NP}}}^P = -H_{-\lambda_2-\lambda_{\text{NP}}}^P$ from parity or explicit calculation. The total left-chiral helicity amplitude can be written as

$$H_{\lambda_2\lambda_W} = H_{\lambda_2\lambda_W}^V - H_{\lambda_2\lambda_W}^A. \quad (12)$$

The angular decay distribution can be written as

$$\frac{d\Gamma}{dq^2 d\cos\theta_l} = \frac{G_F^2 |V_{cb}|^2 |\vec{p}_{B_2}|}{2^9 \pi^3 m_{B_1}^2} \left(1 - \frac{m_l^2}{q^2}\right) L_{\mu\nu} H^{\mu\nu}, \quad (13)$$

where $|\vec{p}_{B_2}| = \sqrt{\lambda(m_{B_1}^2, m_{B_2}^2, q^2)}/2m_{B_1}$ is the momentum of the outgoing baryon B_2 and $\lambda(a, b, c) = a^2 + b^2 + c^2 - 2(ab + bc + ca)$. We denote θ_l as the angle between the daughter baryon B_2 and the lepton three-momentum vector in the q^2 rest frame. Again, $L_{\mu\nu}$ and $H_{\mu\nu}$ are the usual leptonic and hadronic tensors, respectively. The covariant contraction $L_{\mu\nu} H^{\mu\nu}$ can be written in terms of the hadronic and leptonic helicity amplitudes using the helicity techniques. We follow Refs. [21,27,28] and write the differential angular distribution for the three-body $B_1 \rightarrow B_2 l \nu$ decays in the presence of NP as

$$\frac{d\Gamma(B_1 \rightarrow B_2 l \nu)}{dq^2 d\cos\theta_l} = N \left(1 - \frac{m_l^2}{q^2}\right)^2 \left[A_1 + \frac{m_l^2}{q^2} A_2 + 2A_3 + \frac{4m_l}{\sqrt{q^2}} A_4 \right], \quad (14)$$

where

$$\begin{aligned} N &= \frac{G_F^2 |V_{cb}|^2 q^2 |\vec{p}_{B_2}|}{512 \pi^3 m_{B_1}^2}, \\ A_1 &= 2 \sin^2 \theta_l \left(H_{\frac{1}{2}0}^2 + H_{-\frac{1}{2}0}^2 \right) + (1 - \cos \theta_l)^2 H_{\frac{1}{2}1}^2 + (1 + \cos \theta_l)^2 H_{-\frac{1}{2}1}^2, \\ A_2 &= 2 \cos^2 \theta_l \left(H_{\frac{1}{2}0}^2 + H_{-\frac{1}{2}0}^2 \right) + \sin^2 \theta_l \left(H_{\frac{1}{2}1}^2 + H_{-\frac{1}{2}1}^2 \right) + 2 \left(H_{\frac{1}{2}t}^2 + H_{-\frac{1}{2}t}^2 \right) - 4 \cos \theta_l \left(H_{\frac{1}{2}t} H_{\frac{1}{2}0} + H_{-\frac{1}{2}t} H_{-\frac{1}{2}0} \right), \\ A_3 &= (H_{\frac{1}{2}0}^{SP})^2 + (H_{-\frac{1}{2}0}^{SP})^2, \\ A_4 &= -\cos \theta_l \left(H_{\frac{1}{2}0} H_{\frac{1}{2}0}^{SP} + H_{-\frac{1}{2}0} H_{-\frac{1}{2}0}^{SP} \right) + \left(H_{\frac{1}{2}t} H_{\frac{1}{2}0}^{SP} + H_{-\frac{1}{2}t} H_{-\frac{1}{2}0}^{SP} \right). \end{aligned} \quad (15)$$

After integrating the angular distribution of Eq. (14) over $\cos \theta_l$, one obtains

$$\frac{d\Gamma(B_1 \rightarrow B_2 l \nu)}{dq^2} = \frac{8N}{3} \left(1 - \frac{m_l^2}{q^2}\right)^2 \left[B_1 + \frac{m_l^2}{2q^2} B_2 + \frac{3}{2} B_3 + \frac{3m_l}{\sqrt{q^2}} B_4 \right], \quad (16)$$

where

$$\begin{aligned} B_1 &= H_{\frac{1}{2}0}^2 + H_{-\frac{1}{2}0}^2 + H_{\frac{1}{2}1}^2 + H_{-\frac{1}{2}1}^2, \\ B_2 &= H_{\frac{1}{2}0}^2 + H_{-\frac{1}{2}0}^2 + H_{\frac{1}{2}1}^2 + H_{-\frac{1}{2}1}^2 + 3 \left(H_{\frac{1}{2}t}^2 + H_{-\frac{1}{2}t}^2 \right), \\ B_3 &= (H_{\frac{1}{2}0}^{SP})^2 + (H_{-\frac{1}{2}0}^{SP})^2, \\ B_4 &= H_{\frac{1}{2}t} H_{\frac{1}{2}0}^{SP} + H_{-\frac{1}{2}t} H_{-\frac{1}{2}0}^{SP}. \end{aligned} \quad (17)$$

The ratio of branching ratios R_{Ξ_c} is defined as

$$R_{\Xi_c} = \frac{\mathcal{B}(\Xi_b \rightarrow \Xi_c \tau^- \bar{\nu}_\tau)}{\mathcal{B}(\Xi_b \rightarrow \Xi_c l^- \bar{\nu}_l)}, \quad (18)$$

where l is either an electron or a muon. We have also defined several q^2 -dependent observables, such as the differential branching fractions $\text{DBR}(q^2)$, the ratio of branching fractions $R(q^2)$, the forward-backward asymmetries $A_{\text{FB}}^l(q^2)$, the convexity parameter $C_F^l(q^2)$, and the longitudinal polarization fraction of the lepton $P_l(q^2)$ for the baryonic $\Xi_b \rightarrow \Xi_c l \nu$ decay mode. These are

$$\begin{aligned}
\text{DBR}(q^2) &= \left(\frac{d\Gamma}{dq^2} \right) / \Gamma_{\text{tot}}, \\
R(q^2) &= \frac{\text{DBR}(q^2)(B_1 \rightarrow B_2 \tau \nu)}{\text{DBR}(q^2)(B_1 \rightarrow B_2 l \nu)}, \\
P_l(q^2) &= \frac{d\Gamma(+)/dq^2 - d\Gamma(-)/dq^2}{d\Gamma(+)/dq^2 + d\Gamma(-)/dq^2}, \\
A_{FB}^l(q^2) &= \left\{ \left(\int_{-1}^0 - \int_0^1 \right) d \cos \theta_l \frac{d\Gamma}{dq^2 d \cos \theta_l} \right\} / \frac{d\Gamma}{dq^2}, \\
C_F^l(q^2) &= \frac{1}{\mathcal{H}_{\text{tot}}} \frac{d^2 W(\theta)}{d(\cos \theta)^2}, \quad (19)
\end{aligned}$$

where $d\Gamma(+)/dq^2$ and $d\Gamma(-)/dq^2$ denote the differential branching ratios of positive- and negative-helicity leptons, respectively. Again,

$$\begin{aligned}
W(\theta) &= \frac{3}{8} \left[A_1 + \frac{m_l^2}{q^2} A_2 + 2A_3 + \frac{4m_l}{\sqrt{q^2}} A_4 \right], \\
\mathcal{H}_{\text{tot}} &= \int d(\cos \theta) W(\theta), \\
\frac{d^2 W(\theta)}{d(\cos \theta)^2} &= \frac{3}{4} \left(1 - \frac{m_l^2}{q^2} \right) \left[H_{\frac{1}{2}1}^2 + H_{-\frac{1}{2}-1}^2 - 2 \left(H_{\frac{1}{2}0}^2 + H_{-\frac{1}{2}0}^2 \right) \right]. \quad (20)
\end{aligned}$$

We also give predictions for the average values of the forward-backward asymmetry of the charged lepton $\langle A_{FB}^l \rangle$, the convexity parameter $\langle C_F^l \rangle$, and the longitudinal polarization of the lepton $\langle P_l \rangle$ which are calculated by separately integrating the numerators and denominators over q^2 .

III. RESULTS AND DISCUSSION

For definiteness, we first present all the inputs that are relevant for our numerical computation. For the quark, lepton, and baryon masses we use $m_b(m_b) = 4.18$ GeV, $m_c(m_b) = 0.91$ GeV, $m_e = 0.510998928 \times 10^{-3}$ GeV, $m_\mu = 0.1056583715$ GeV, $m_\tau = 1.77682$ GeV, $m_{\Xi_b} = 5.7919$ GeV, and $m_{\Xi_c} = 2.46787$ GeV [29]. For the mean lifetime of the Ξ_b baryon we use $\tau_{\Xi_b} = 1.479 \times 10^{-12}$ s [29]. For the CKM matrix element $|V_{cb}|$ we use the value $|V_{cb}| = (40.9 \pm 1.1) \times 10^{-3}$ [29]. The relevant parameters for the form factor calculation taken from Ref. [8] are reported in Table I. We use $\pm 10\%$ uncertainty in each of these parameters. We also report the most important experimental input parameters R_D and R_{D^*} in Table II.

TABLE I. Parameters for the Isgur-Wise functions of $\Xi_b \rightarrow \Xi_c$ form factors taken from Ref. [8].

$\bar{\Lambda}(\text{GeV})$	ρ_ζ^2	c_ζ	ρ_χ^2	c_χ
0.970	2.27	3.87	0.045	0.036

TABLE II. Current status of R_D and R_{D^*} [1].

Experiments	R_{D^*}	R_D
BABAR	$0.332 \pm 0.024 \pm 0.018$	$0.440 \pm 0.058 \pm 0.042$
BELLE	$0.293 \pm 0.038 \pm 0.015$	$0.375 \pm 0.064 \pm 0.026$
BELLE	$0.302 \pm 0.030 \pm 0.011$	
LHCb	$0.336 \pm 0.027 \pm 0.030$	
BELLE	$0.270 \pm 0.035^{+0.028}_{-0.025}$	
LHCb	$0.285 \pm 0.019 \pm 0.029$	
AVERAGE	$0.304 \pm 0.013 \pm 0.007$	$0.407 \pm 0.039 \pm 0.024$

We use the average values of R_D and R_{D^*} for our analysis. For the errors, we add the statistical and systematic uncertainties in quadrature.

There are two major sources of uncertainty in the calculation of the decay amplitudes. It may come from either not so well-known input parameters such as the CKM matrix elements or hadronic input parameters such as form factors and decay constants. In order to gauge the effect of these uncertainties on various observables, we use a random number generator and perform a random scan over all the theoretical input parameters—such as the CKM matrix elements, form factors, and decay constants—within 1σ of their central values. The SM central values and the corresponding 1σ ranges of all the observables for the $\Xi_b \rightarrow \Xi_c l \nu$ decays are presented in Table III. We notice that there are considerable changes while going from the e to the τ mode, including even a sign change in the forward-backward asymmetry parameter $\langle A_{FB}^l \rangle$. The central values reported in Table III are obtained using the central values of all the input parameters, whereas to find the 1σ range of all the observables, we vary all the input parameters such as CKM matrix elements, the hadronic form factors, and the decay constants within 1σ of their central values. However, we do not include the uncertainties coming from the quark mass, lepton mass, baryon mass, and the mean lifetime as these are not important for our analysis.

In Fig. 1, we show the q^2 dependence of $A_{FB}^l(q^2)$, $C_F^l(q^2)$, and $P_l(q^2)$ within the SM for the τ and e modes. We observe that the q^2 behavior of all the observables for the e mode is quite different from the τ mode. The forward-backward asymmetry parameter $A_{FB}^l(q^2)$ approaches zero

TABLE III. Predictions of various observables within the SM.

Observables	Central value	1σ range
$\mathcal{B}(\Xi_b \rightarrow \Xi_c e \nu)\%$	9.22	[5.62, 14.47]
$\mathcal{B}(\Xi_b \rightarrow \Xi_c \tau \nu)\%$	2.35	[1.84, 2.99]
R_{Ξ_c}	0.255	[0.205, 0.330]
$\langle A_{FB}^e \rangle$	0.163	[0.146, 0.185]
$\langle A_{FB}^\tau \rangle$	-0.042	[-0.027, -0.056]
$\langle C_F^e \rangle$	-0.697	[-0.572, -0.781]
$\langle C_F^\tau \rangle$	-0.103	[-0.100, -0.106]
$\langle P_e \rangle$	-1.00	-1.00
$\langle P_\tau \rangle$	-0.317	[-0.298, -0.338]

FIG. 1. q^2 dependence of $A_{FB}^l(q^2)$, $C_F^l(q^2)$, and $P_l(q^2)$ within the SM. Blue and green lines correspond to the τ and e modes, respectively.

at zero recoil for both the e and τ modes. We observe that although $A_{FB}^l(q^2)$ remains positive for the e mode, it becomes negative for the τ mode below $q^2 \equiv 8.0 \text{ GeV}^2$. We observe a zero crossing in the $A_{FB}^l(q^2)$ parameter for the τ mode. Similarly, at zero recoil the $C_F^l(q^2)$ parameter approaches zero for both the e and τ modes. However, at maximum recoil $C_F^l(q^2)$ becomes zero for the τ mode, whereas it becomes large and negative for the e mode. Again, the convexity parameter remains very small in the entire q^2 region for the τ mode. The longitudinal polarization fraction of the charged lepton $P_l(q^2)$ is -1 in the entire q^2 region for the e mode. For the τ mode, we observe a zero crossing in the $P_l(q^2)$ parameter at $q^2 \equiv 4.6 \text{ GeV}^2$, below which it becomes positive.

Now we proceed to discuss various NP scenarios. We want to see the effect of various NP couplings in a model-independent way. In the first scenario, we assume that NP is coming from couplings associated with new vector interactions, i.e., from V_L and V_R only. We vary V_L and V_R while keeping $S_{L,R} = 0$. In order to determine the allowed NP parameter space, we impose the 1σ constraint coming from the measured values of the ratio of branching ratios R_D and R_{D^*} . The allowed ranges in V_L and V_R that satisfy the 1σ experimental constraint are shown in the left panel of Fig. 2. In the right panel, we show the allowed ranges in $\Gamma(B_c \rightarrow \tau\nu)$ and $\mathcal{B}(\Xi_b \rightarrow \Xi_c \tau\nu)$ obtained in this NP scenario. We see that $\Gamma(B_c \rightarrow \tau\nu)$ obtained in this scenario is

consistent with $\Gamma(B_c \rightarrow \tau\nu) \leq 30\%$ of the measured experimental width, Γ_{B_c} . The corresponding ranges of all the observables are listed in Table IV. We see a significant deviation from the SM prediction. Depending on the NP couplings V_L and V_R , the values of the observables can be either smaller or larger than the SM prediction.

We wish to look at the effect of the new physics couplings (V_L, V_R) on different observables, such as the differential branching ratio $\text{DBR}(q^2)$, ratio of branching ratios $R(q^2)$, forward-backward asymmetry $A_{FB}^l(q^2)$, convexity parameter $C_F^l(q^2)$, and τ polarization fraction $P_\tau(q^2)$ for the $\Xi_b \rightarrow \Xi_c \tau\nu$ decays. In Fig. 3, we show in blue the allowed SM bands and in green the allowed bands of each observable once the NP couplings V_L and V_R are switched on. It can be seen that once NP is included the deviation from the SM expectation is quite large in the cases of $\text{DBR}(q^2)$, $R(q^2)$, and $A_{FB}^l(q^2)$. However, the deviation is slightly less in the cases of $C_F^l(q^2)$ and $P_\tau(q^2)$. We observe that, depending on the values of NP couplings, there may or may not be a zero crossing in the forward-backward asymmetry parameter $A_{FB}^l(q^2)$. In the case of $P_\tau(q^2)$, the zero-crossing point may shift slightly towards a higher q^2 value than in the SM.

In the second scenario, we assume that NP is coming from new scalar interactions, i.e., from S_L and S_R only. To explore the effect of NP coming from S_L and S_R , we vary S_L and S_R and impose the 1σ constraint coming from the

FIG. 2. Allowed regions of V_L and V_R obtained using the 1σ constraint coming from R_D and R_{D^*} are shown in the left panel and the corresponding ranges in $\Gamma(B_c \rightarrow \tau\nu)$ and $\mathcal{B}(\Xi_b \rightarrow \Xi_c \tau\nu)$ in the presence of these NP couplings are shown in the right panel. The blue horizontal line in the right panel corresponds to $\Gamma(B_c \rightarrow \tau\nu) = 30\%$ of the measured experimental width, Γ_{B_c} .

TABLE IV. Ranges of various observables with V_L and V_R NP couplings from the left panel of Fig. 2.

$\mathcal{B}(\Xi_b \rightarrow \Xi_c \tau \nu)\%$	R_{Ξ_c}	$\langle A_{FB}^\tau \rangle$	$\langle C_F^\tau \rangle$	$\langle P_\tau \rangle$
[2.36, 3.71]	[0.254, 0.410]	[-0.028, -0.346]	[-0.101, -0.113]	[-0.246, -0.333]

measured values of R_D and R_{D^*} . The resulting ranges in S_L and S_R obtained using the 1σ experimental constraint are shown in the left panel of Fig. 4. It should be mentioned that our analysis of new physics effects on $\Xi_b \rightarrow \Xi_c \tau \nu$ decays differs slightly from our previous analysis in Ref. [21]. In Ref. [21], we used the constraints coming from R_D , R_{D^*} , and R_π^l where $R_\pi^l = \frac{\tau_{B_0}}{\tau_{B^-}} \frac{\mathcal{B}(B \rightarrow \tau \nu)}{\mathcal{B}(B \rightarrow \pi l \nu)}$ to determine the allowed NP parameter space. However, in the present analysis we use the constraints coming from R_D and R_{D^*}

alone to determine the allowed NP parameter space. Due to this additional constraint from R_π^l , the allowed (S_L, S_R) NP parameter space in Fig. 3 of Ref. [21] is more constrained than in the left panel of Fig. 4. We notice that the allowed ranges of $\Gamma(B_c \rightarrow \tau \nu)$ obtained in this scenario are rather large—more than 30% of the measured experimental width Γ_{B_c} . It should also be noted that for some values of S_L and S_R , $\Gamma(B_c \rightarrow \tau \nu)$ can also be larger than the experimental width of the B_c meson, which is not acceptable at all. We

FIG. 3. The dependence of the observables $\text{DBR}(q^2)$, $R(q^2)$, $A_{FB}^\tau(q^2)$, $C_F^\tau(q^2)$, and $P_\tau(q^2)$ on V_L and V_R NP couplings. The allowed range in each observable is shown with light green band once the NP couplings (V_L, V_R) are varied within the allowed ranges from the left panel of Fig. 2. The corresponding SM prediction is shown with light blue band.

FIG. 4. Allowed regions of S_L and S_R obtained using the 1σ constraint coming from R_D and R_{D^*} are shown in the left panel and the corresponding ranges in $\Gamma(B_c \rightarrow \tau \nu)$ and $\mathcal{B}(\Xi_b \rightarrow \Xi_c \tau \nu)$ in the presence of these NP couplings are shown in the right panel. The blue horizontal line in the right panel corresponds to $\Gamma(B_c \rightarrow \tau \nu) = 30\%$ of Γ_{B_c} , whereas the green horizontal line denotes the measured experimental width, Γ_{B_c} .

TABLE V. Ranges of various observables with S_L and S_R NP couplings from the left panel of Fig. 4.

$\mathcal{B}(\Xi_b \rightarrow \Xi_c \tau \nu)$	R_{Ξ_c}	$\langle A_{FB}^\tau \rangle$	$\langle C_F^\tau \rangle$	$\langle P_\tau \rangle$
[1.94, 5.42]	[0.200, 0.616]	[-0.179, 0.226]	[-0.052, -0.111]	[-0.371, 0.338]

illustrate these points in the right panel of Fig. 4. Thus, any explanation of R_D and R_{D^*} anomalies with S_L and S_R NP couplings is disfavored by the constraint coming from the experimental width of the B_c meson. Nevertheless, we report in Table V the allowed ranges of all the observables in this scenario. We see a significant deviation of all the observables from the SM prediction. It should be noted that the deviation observed in this scenario is more pronounced than the deviation observed with V_L and V_R NP couplings.

We want to see the effect of these NP couplings on various q^2 -dependent observables. In Fig. 5, we show how the observables $\text{DBR}(q^2)$, $R(q^2)$, $A_{FB}^\tau(q^2)$, $C_F^\tau(q^2)$, and $P_\tau(q^2)$ behave as a function of q^2 with and without S_L and S_R NP couplings. The light blue band corresponds to the SM range, whereas the light green band corresponds to the range of the observables with S_L and S_R NP couplings. The deviations from the SM expectation are prominent for each observable. It should be mentioned that the deviation observed in this scenario is more pronounced than the deviation observed with V_L and V_R NP couplings. We see that the $\text{DBR}(q^2)$ for the $\Xi_b \rightarrow \Xi_c \tau \nu$ decay mode can increase by more than 100% at the peak of its distribution. Similar conclusions can be made for the ratio of branching ratios as well. Again, depending on the values of S_L and S_R NP couplings, the zero crossing points of $A_{FB}^\tau(q^2)$ and $P_\tau(q^2)$ can be quite different from the SM prediction. Moreover, it is evident

from Fig. 5 that depending on S_L and S_R there may not be any zero crossing for $A_{FB}^\tau(q^2)$ and $P_\tau(q^2)$.

IV. SUMMARY AND CONCLUSION

Lepton flavor universality violation has been reported in various semileptonic B meson decays. Tensions between SM predictions and experiments exist in various semileptonic B meson decays mediated via $b \rightarrow c$ charged-current interactions and $b \rightarrow s \bar{l} l$ neutral-current interactions. The study of $\Xi_b \rightarrow \Xi_c \tau \nu$ decays is important mainly for two reasons. First, it can act as a complimentary decay channel to $B \rightarrow (D, D^*) \tau \nu$ decays mediated via $b \rightarrow c$ charged-current interactions and, in principle, can provide complementary constraints on possible origins of the anomalies present in the ratios R_D and R_{D^*} . Second, the precise determination of the branching fractions of this decay mode will allow an accurate determination of the CKM matrix element $|V_{cb}|$ with less theoretical uncertainty.

We have used the helicity formalism to study baryonic $\Xi_b \rightarrow \Xi_c l \nu$ decays within the context of an effective Lagrangian in the presence of NP. We have defined various observables and provided predictions using form factors obtained in the relativistic quark model. We have given the first predictions of various observables for this decay mode, such as R_{Ξ_c} , A_{FB}^l , P_l , and C_F^l . We also looked at possible

FIG. 5. The dependence of the observables $\text{DBR}(q^2)$, $R(q^2)$, $A_{FB}^\tau(q^2)$, $C_F^\tau(q^2)$, and $P_\tau(q^2)$ on S_L and S_R NP couplings. The allowed range in each observable is shown with light green band once the NP couplings (S_L, S_R) are varied within the allowed ranges shown in the left panel of Fig. 4. The corresponding SM prediction is shown with light blue band.

NP effects on various observables for this decay mode. Let us now summarize our main results.

We first reported the central values and the 1σ ranges of all the observables for the $\Xi_b \rightarrow \Xi_c l \nu$ decays within the SM. The SM branching ratio of $\Xi_b \rightarrow \Xi_c l \nu$ decays is of the order of 10^{-1} . We observed that the integrated quantities—forward-backward asymmetry $\langle A_{FB}^l \rangle$, longitudinal polarization fraction of the charged lepton $\langle P_l \rangle$, and the convexity parameter $\langle C_F^l \rangle$ —change considerably when going from the e to the τ mode. There is even a sign change in the case of the forward-backward asymmetry parameter $\langle A_{FB}^l \rangle$. We also studied the q^2 dependence of various observables within the SM for both the τ and e modes. We observed considerable changes in each of the observables when going from the e to the τ mode.

For the NP analysis, we included vector and scalar NP interactions and explored two different NP scenarios. In the first scenario, we considered that only new vector NP interactions characterized by V_L and V_R NP couplings contribute to the decay mode. In the second scenario, we assumed that NP comes from new scalar interactions, i.e., from S_L and S_R NP couplings only. The allowed ranges in the NP couplings were obtained by using the 1σ constraint coming from the measured values of R_D and R_{D^*} . We also studied the effect of these NP couplings on various

q^2 -dependent observables, such as $\text{DBR}(q^2)$, $R(q^2)$, $A_{FB}^\tau(q^2)$, $C_F^\tau(q^2)$, and $P_\tau(q^2)$. We found significant deviations from the SM prediction once the NP couplings were included. However, the deviation from the SM prediction is more pronounced in the case of scalar NP interactions S_L and S_R . It should be mentioned that $\mathcal{B}(B_c \rightarrow \tau \nu)$ puts a severe constraint on the S_L and S_R NP couplings. Although S_L and S_R NP couplings can simultaneously explain the anomalies present in R_D and R_{D^*} , they do not satisfy the $\mathcal{B}(B_c \rightarrow \tau \nu) \leq 30\%$ constraint. Hence, the conclusion is that the B_c lifetime makes highly implausible any explanation of R_D and R_{D^*} anomalies with S_L and S_R NP couplings. With V_L and V_R NP couplings, the allowed range obtained for $\mathcal{B}(B_c \rightarrow \tau \nu)$ is consistent with the $\mathcal{B}(B_c \rightarrow \tau \nu) \leq 30\%$ obtained in the SM.

Although there is a hint of NP in the meson sector, NP is not yet established. The study of $\Xi_b \rightarrow \Xi_c l \nu$ decays both theoretically and experimentally is well motivated because of the longstanding anomalies present in R_D and R_{D^*} . It would be interesting to find a similar hint of NP in the semileptonic baryonic decays as well. At the same time, a precise measurement of $\mathcal{B}(\Xi_b \rightarrow \Xi_c l \nu)$ and a precise determination of $\Xi_b \rightarrow \Xi_c$ transition form factors will allow an accurate determination of the CKM matrix element $|V_{cb}|$.

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