Charmful two-body anti-triplet $b$-baryon decays

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Abstract

We study the charmful decays of the two-body $B_b \to B_n M_c$ decays, where $B_b$ represents the anti-triplet of $(\Lambda_b, \Xi^0_b, \Xi^-_b)$, $B_n$ stands for the baryon octet and $M_c$ denotes as the charmed meson of $D^{(*)}_{(s)}$, $\eta_c$ and $J/\psi$. Explicitly, we predict that $\mathcal{B}(\Lambda_b \to D^-_s p) = (1.8 \pm 0.3) \times 10^{-5}$, which is within the measured upper bound of $\mathcal{B}(\Lambda_b \to D^-_s p) < 4.8(5.3) \times 10^{-4}$ at 90% (95%) C.L., and reproduce $\mathcal{B}(\Lambda_b \to J/\psi \Lambda) = (3.3 \pm 2.0) \times 10^{-4}$ and $\mathcal{B}(\Xi^-_b \to J/\psi \Xi^-) = (5.1 \pm 3.2) \times 10^{-4}$ in agreement with the data. Moreover, we find that $\mathcal{B}(\Lambda_b \to \Lambda \eta_c) = (1.5 \pm 0.9) \times 10^{-4}$, $\mathcal{B}(\Xi^-_b \to \Xi^- \eta_c) = (2.4 \pm 1.5) \times 10^{-4}$ and $\mathcal{B}(\Xi^0_b \to \Xi^0 \eta_c, \Xi^0 J/\psi) = (2.3 \pm 1.4, 4.9 \pm 3.0) \times 10^{-4}$, which are accessible to the experiments at the LHCb.
I. INTRODUCTION

The two-body decays of $\Lambda_b \to \Lambda_c^+ K^-, \Lambda_c^+ \pi^-, \Lambda_c^+ D^-$, and $\Lambda_c^+ D_s^-$ can be viewed as through the $\Lambda_b \to \Lambda_c$ transition along with the recoiled mesons $K^-, \pi^-, D_s^-$ and $D^-$, respectively, such that one may use the factorization ansatz to get the fractions of the branching ratios as

\[
\mathcal{R}_{K/\pi} = \frac{\mathcal{B}(\Lambda_b \to \Lambda_c^+ K^-)}{\mathcal{B}(\Lambda_b \to \Lambda_c^+ \pi^-)} \approx \frac{|V_{us}|^2 f_K^2}{|V_{ud}|^2 f_\pi^2} = 0.073,
\]

\[
\mathcal{R}_{D/D_s} = \frac{\mathcal{B}(\Lambda_b \to \Lambda_c^+ D^-)}{\mathcal{B}(\Lambda_b \to \Lambda_c^+ D_s^-)} \approx \frac{|V_{us}|^2 f_D^2}{|V_{us}|^2 f_{D_s}^2} = 0.034,
\]

which are in agreement with the data, given by $[1, 2]$.

\[
\mathcal{R}_{K/\pi} = 0.0731 \pm 0.0016 \pm 0.0016, \quad \mathcal{R}_{D/D_s} = 0.042 \pm 0.003 \pm 0.003.
\]

In the same picture, the measured $\mathcal{B}(\Lambda_b \to pK^-, p\pi)$ can be also explained $[3, 4]$. In addition, the direct CP violating asymmetry of $\Lambda_b \to pK^{*-}$ is predicted as large as 20% $[5]$.

On the other hand, the branching ratios of $\Lambda_b \to D_s^- p$, $\Lambda_b \to J/\psi \Lambda$ and $\Xi_b^- \to J/\psi \Xi^-$ are shown as $[4, 6]$:

\[
\mathcal{B}(\Lambda_b \to D_s^- p) = (2.7 \pm 1.4 \pm 0.2 \pm 0.7 \pm 0.1 \pm 0.1) \times 10^{-4} \text{ or } 2.4 \pm 1.5 \times 10^{-4} \text{ at the } 90\% \text{ (95\%)} \text{ C.L.},
\]

\[
\mathcal{B}(\Lambda_b \to J/\psi \Lambda) = (3.0 \pm 1.1) \times 10^{-4},
\]

\[
\mathcal{B}(\Xi_b^- \to J/\psi \Xi^-) = (2.0 \pm 0.9) \times 10^{-4},
\]

with $\mathcal{B}(\Lambda_b \to J/\psi \Lambda)$ and $\mathcal{B}(\Xi_b^- \to J/\psi \Xi^-)$ converted from the partial observations of $\mathcal{B}(\Lambda_b \to J/\psi \Lambda)f_{\Lambda_b} = (5.8 \pm 0.8) \times 10^{-5}$ and $\mathcal{B}(\Xi_b^- \to J/\psi \Xi^-) f_{\Xi_b^-} = (1.02^{+0.20}_{-0.21}) \times 10^{-5}$, where $f_{\Lambda_b} = 0.175 \pm 0.106$ and $f_{\Xi_b^-} = 0.019 \pm 0.013$ are the fragmentation fractions of the $b$ quark to $b$-baryons of $\Lambda_b$ and $\Xi_b$ $[7]$, respectively. Nonetheless, for these $B_b \to B_n M_c$ decays in Eq. (3), the theoretical understanding is still lacking. Since the factorization approach is expected to be reliable in studying the branching ratios of $B_b \to B_n M_c$, in this report, we shall systematically analyze the branching ratios for all possible $B_b \to B_n M_c$ decays, and compare them with the experimental data at the $B$-factories, as well as the LHCb, where $B_b$, $B_n$ and $M_c$ correspond to the anti-triplet $b$-baryon of $(\Lambda_b, \Xi_b^0, \Xi_b^-)$, baryon octet and charmed meson, respectively.
II. FORMALISM

As the studies in Refs. [8–14], based on the factorization approach, the amplitudes for the two-body charmful \( B_b \rightarrow B_n M_c \) decays are presented in terms of the decaying process of the \( B_b \rightarrow B_n \) transition with the recoiled charmed meson \( M_c \). From Fig. 1(a), the amplitudes of \( B_b \rightarrow B_n M_c \) via the quark-level \( b \rightarrow u\bar{c}q \) transition are factorized as

\[
A_1(B_b \rightarrow B_n M_c) = G_F \sqrt{2} V_{ub} V_{cq}^* a_1 \langle M_c | \bar{q}\gamma^\mu (1 - \gamma_5) c|0\rangle \langle B_n | \bar{u}\gamma^\mu (1 - \gamma_5) b|B_b \rangle , \tag{4}
\]

where \( G_F \) is the Fermi constant, \( V_{ub,cq} \) are the CKM matrix elements, while the explicit decay modes are

\[
\Lambda_b \rightarrow pM_c, \quad \Xi_b^- \rightarrow \Lambda(\Sigma^0)M_c, \quad \Xi_b^0 \rightarrow \Sigma^+ M_c \tag{5}
\]

with \( q = d(s) \) for \( M_c = D^{(*)-}(D_s^{(*)-}) \). On the other hand, the amplitudes via the quark-level \( b \rightarrow c\bar{u}q \) \((b \rightarrow c\bar{u}q)\) transition in Fig. 1(b) can be written as

\[
A_2(B_b \rightarrow B_n M_c) = G_F \sqrt{2} V_{cb} V_{q_1q}^* a_2 \langle M_c | \bar{c}\gamma^\mu (1 - \gamma_5) q_1|0\rangle \langle B_n | \bar{q}\gamma^\mu (1 - \gamma_5) b|B_b \rangle , \tag{6}
\]

with \( q_1 = u \) for \( M_c = D^{(*)0} \) and \( q_1 = c \) for \( M_c = \eta_c \) and \( J/\psi \), where the decays of \( B_b \rightarrow B_n M_c \) are

\[
\begin{align*}
\Lambda_b \rightarrow n M_c, \quad \Xi_b^- \rightarrow \Sigma^- M_c, \quad \Xi_b^0 \rightarrow \Lambda(\Sigma^0)M_c & \quad \text{for} \ q_2 = d, \\
\Lambda_b \rightarrow \Lambda(\Sigma^0) M_c, \quad \Xi_b^- \rightarrow \Xi^- M_c, \quad \Xi_b^0 \rightarrow \Xi^0 M_c & \quad \text{for} \ q_2 = s. \tag{7}
\end{align*}
\]

In this study, we will exclude the study of \( \Lambda_b \rightarrow n M_c \) due to the elusive neutron in the \( B \)-factories. The amplitudes \( A_{1,2} \) via the \( W \)-boson exchange diagrams are led to be the color-allowed and color-suppressed processes. The parameters \( a_1 \) and \( a_2 \) in Eqs. (4) and (6)
are presented as \[15,16\]

\[
a_1 = c_1^{\text{eff}} + \frac{c_2^{\text{eff}}}{N_c}, \quad a_2 = c_2^{\text{eff}} + \frac{c_1^{\text{eff}}}{N_c},
\]

with the effective Wilson coefficients \((c_1^{\text{eff}}, c_2^{\text{eff}}) = (1.168, -0.365)\), respectively, where the color number \(N_c\) should be taken as a floating number from 2 \(\rightarrow\) \(\infty\) to account for the non-factorizable effects in the generalized factorization instead of \(N_c = 3\). The matrix elements for \(P_c = (\eta_c, D)\) and \(V_c = J(\psi, D^*)\) productions read

\[
\langle P_c | A_{\mu}^c | 0 \rangle = -i f_{P_c} q_{\mu}, \\
\langle V_c | V_{\mu}^c | 0 \rangle = m_{V_c} f_{V_c} \varepsilon_{\mu}^*,
\]

with \(V_{\mu}^c(A_{\mu}^c) = \bar{q} \gamma_{\mu}(\gamma_5)c\) or \(\bar{c} \gamma_{\mu}(\gamma_5)q_1\), where \(q_{\mu}\) and \(\varepsilon_{\mu}^*\) are the four-momentum and polarization, respectively. Those of the \(B_b \rightarrow B_n\) baryon transition in Eq. (11) have the general forms:

\[
\langle B_n | \bar{q} \gamma_{\mu} b | B_b \rangle = \bar{u}_{B_n} \left[ f_1 \gamma_{\mu} + \frac{f_2}{m_{B_b}} i \sigma_{\mu\nu} q^\nu + \frac{f_3}{m_{B_b}^2} q_{\mu} \right] u_{B_b},
\]

\[
\langle B_n | \bar{q} \gamma_{\mu} \gamma_5 b | B_b \rangle = \bar{u}_{B_n} \left[ g_1 \gamma_{\mu} + \frac{g_2}{m_{B_b}} i \sigma_{\mu\nu} q^\nu + \frac{g_3}{m_{B_b}^2} q_{\mu} \right] \gamma_5 u_{B_b},
\]

where \(f_j (g_j) (j = 1, 2, 3)\) are the form factors. We are able to relate the different \(B_b \rightarrow B_n\) transition form factors based on the \(SU(3)\) flavor and \(SU(2)\) spin symmetries, which have been used to connect the space-like \(B_n \rightarrow B_n'\) transition form factors in the neutron decays \[17\], and the time-like \(0 \rightarrow B_n B_n'\) baryonic form factors as well as the \(B \rightarrow B_n B_n'\) transition form factors in the baryonic \(B\) decays \[18\,22\]. Specifically, \(V_{\mu} = \bar{q} \gamma_{\mu} b\) and \(A_{\mu} = \bar{c} \gamma_{\mu} \gamma_5 b\) as the two currents in Eq. (10) can be combined as the right-handed chiral current, that is, \(J_{\mu,R}^q = (V_{\mu} + A_{\mu})/2\). Consequently, we have \[17\]:

\[
\langle B_n,^{\uparrow+\downarrow} | J_{\mu,R}^q | B_{b,^{\uparrow+\downarrow}} \rangle = \bar{u}_{B_n} \left[ \gamma_{\mu} \frac{1 + \gamma_5}{2} G^\dagger(q^2) + \gamma_{\mu} \frac{1 - \gamma_5}{2} G^\dagger(q^2) \right] u_{B_b},
\]

where the baryon helicity states \(|B_{n(b),^{\uparrow+\downarrow}}\rangle \equiv |B_{n(b),^{\uparrow}}\rangle + |B_{n(b),^{\downarrow}}\rangle\) are regarded as the baryon chiral states \(|B_{n(b),R+L}\rangle\) at the large momentum transfer, while \(G^\dagger(q^2)\) and \(G^\dagger(q^2)\) are the right-handed and left-handed form factors, defined by

\[
G^\dagger(q^2) = e^\dagger_\parallel G_\parallel(q^2) + e^\dagger_\perp G_\perp(q^2), \quad G^\dagger(q^2) = e^\dagger_\parallel G_\parallel(q^2) + e^\dagger_\perp G_\perp(q^2),
\]
with the constants \( e^+_\parallel\) and \( e^+_\perp\) to sum over the chiral charges via the \( B_b \to B_n \) transition, given by

\[
e^+_\parallel = \langle B_{n,\parallel}|Q_\parallel|B_{b,\downarrow}\rangle, \quad e^+_\perp = \langle B_{n,\downarrow}|Q_\parallel|B_{b,\uparrow}\rangle,
\]
\[
e^-\parallel = \langle B_{n,\downarrow}|Q_\parallel|B_{b,\downarrow}\rangle, \quad e^-\perp = \langle B_{n,\uparrow}|Q_\parallel|B_{b,\downarrow}\rangle.
\]

(13)

Note that \( Q_\parallel(i) = \sum Q_\parallel(i) \) with \( i = 1, 2, 3 \) as the the chiral charge operators are from \( Q_R^q \equiv J_{0,R}^q b_R \), converting the \( b \) quark in \( |B_{b,\uparrow}\rangle \) into the \( q \) one, while the converted \( q \) quark can be parallel or antiparallel to the \( B_b \)’s helicity, denoted as the subscript (\( \parallel \) or \( \perp \)). By comparing Eq. (10) with Eqs. (11), (12), and (13), we obtain

\[
f_1 = (e^+_\parallel + e^+_\perp)G_\parallel + (e^-\parallel + e^-\perp)G_\perp,
\]
\[
g_1 = (e^+_\parallel - e^-\parallel)G_\parallel + (e^-\perp - e^+\perp)G_\perp,
\]

(14)

with \( f_{2,3} = 0 \) and \( g_{2,3} = 0 \) due to the helicity conservation, as those derived in Refs. [3, 10, 23]. It is interesting to see that, as the helicity-flip terms, the theoretical calculations from the loop contributions to \( f_{2,3} \) (\( g_{2,3} \)) indeed result in the values to be one order of magnitude smaller than that of \( f_1(g_1) \), which can be safely neglected. In the double-pole momentum dependences, \( f_1 \) and \( g_1 \) can be given as

\[
f_1(q^2) = \frac{f_1(0)}{(1 - q^2/m_{B_b}^2)^2}, \quad g_1(q^2) = \frac{g_1(0)}{(1 - q^2/m_{B_b}^2)^2},
\]

(15)

such that it is reasonable to parameterize the chiral form factors to be \( (1 - q^2/m_{B_b}^2)^2 G_\parallel = C_\parallel \). Subsequently, from

\[
(e^\uparrow_\parallel, e^\downarrow_\parallel, e^\uparrow_\perp, e^\downarrow_\perp) = (-\sqrt{3}/2, 0, 0, 0) \quad \text{for} \quad \langle p|J^u_{\mu,R}|\Lambda_b\rangle,
\]
\[
(e^\uparrow_\parallel, e^\downarrow_\parallel, e^\uparrow_\perp, e^\downarrow_\perp) = (1, 0, 0, 0) \quad \text{for} \quad \langle \Lambda|J^a_{\mu,R}|\Lambda_b\rangle,
\]
\[
(e^\uparrow_\parallel, e^\downarrow_\parallel, e^\uparrow_\perp, e^\downarrow_\perp) = (0, 0, 0, 0) \quad \text{for} \quad \langle \Sigma^0|J^s_{\mu,R}|\Lambda_b\rangle,
\]

(16)

we get \( f_1(0) = g_1(0) = -\sqrt{3}/2 C_\parallel \) for \( \langle p|\bar{u}\gamma_\mu(\gamma_5)b|\Lambda_b\rangle \), \( f_1(0) = g_1(0) = C_\parallel \) for \( \langle \Lambda|\bar{s}\gamma_\mu(\gamma_5)b|\Lambda_b\rangle \), and \( f_1(0) = g_1(0) = 0 \) for \( \langle \Sigma^0|\bar{s}\gamma_\mu(\gamma_5)b|\Lambda_b\rangle \), similar to the results based on the heavy-quark and large-energy symmetries in Ref. [23] for the \( \Lambda_b \to (p, \Lambda, \Sigma) \) transitions. When we further extend the study to the anti-triplet \( b \)-baryons: \( (\Xi^-_b, \Xi^0_b, \Lambda^0_b) \) shown in Table II we find that the relation of \( f_1 = g_1 \) is uniquely determined for the anti-triplet \( b \)-baryon transitions.
TABLE I. Relations between the transition matrix elements.

| Transition | Relation |
|------------|----------|
| $\langle B_n| \bar{q}b |B_b \rangle$ | $f_1(0) = g_1(0)$ |
| $\langle p| \bar{u}b |\Lambda_b \rangle$ | $-\sqrt{\frac{3}{2}} C_{||}$ |
| $\langle \Lambda| \bar{u}b |\Xi_b^- \rangle$ | $\frac{1}{2} C_{||}$ |
| $\langle \Sigma^0| \bar{u}b |\Xi_b^- \rangle$ | $-\sqrt{\frac{3}{2}} C_{||}$ |
| $\langle \Sigma^+| \bar{u}b |\Xi_b^0 \rangle$ | $-\sqrt{\frac{3}{2}} C_{||}$ |
| $\langle \Sigma^-| \bar{d}b |\Xi_b^- \rangle$ | $\sqrt{\frac{3}{2}} C_{||}$ |
| $\langle \Lambda| \bar{d}b |\Xi_b^0 \rangle$ | $-\frac{1}{2} C_{||}$ |
| $\langle \Sigma^0| \bar{d}b |\Xi_b^0 \rangle$ | $\sqrt{\frac{3}{2}} C_{||}$ |
| $\langle \Lambda| \bar{s}b |\Lambda_b \rangle$ | $C_{||}$ |
| $\langle \Sigma^0| \bar{s}b |\Lambda_b \rangle$ | $0$ |
| $\langle \Xi^-| \bar{s}b |\Xi_b^- \rangle$ | $\sqrt{\frac{3}{2}} C_{||}$ |
| $\langle \Xi^0| \bar{s}b |\Xi_b^0 \rangle$ | $-\sqrt{\frac{3}{2}} C_{||}$ |

III. NUMERICAL RESULTS AND DISCUSSIONS

For the numerical analysis, the CKM matrix elements in the Wolfenstein parameterization taken from the PDG [4] are given by

$$(V_{ub}, V_{cb}) = (A \lambda^3 (\rho - i \eta), A \lambda^2),$$

$$(V_{cd} = -V_{us}, V_{cs} = V_{ud} = -\lambda, 1 - \lambda^2/2),$$

with $(\lambda, A, \rho, \eta) = (0.225, 0.814, 0.120 \pm 0.022, 0.362 \pm 0.013)$. The meson decay constants are adopted as $(f_{\eta}, f_{J/\psi}) = (387 \pm 7, 418 \pm 9)$ MeV [24], $(f_D, f_{D_s}) = (204.6 \pm 5.0, 257.5 \pm 4.6)$ MeV [4], and $(f_{D^*}, f_{D_s^*}) = (252.2 \pm 22.7, 305.5 \pm 27.3)$ MeV [25]. As given in Ref. [3] to explain the branching ratios and CP violating asymmetries of $\Lambda_b \to p(K^-, \pi^-)$, we have $|\sqrt{3/2} C_{||}| = 0.136 \pm 0.009$ for $\langle p| \bar{u} \gamma_\mu (\gamma_5)b|\Lambda_b \rangle$, which is consistent with the value of $0.14 \pm 0.03$ in the light-cone sum rules [23] and the theoretical calculations in Refs. [8, 10].

With $B(\Lambda_b \to J/\psi \Lambda)$ and $B(\Xi_b^- \to J/\psi \Xi^-)$ in Eq. (8) as the experimental inputs, we can estimate the non-factorizable effects by deviating the color number $N_c = 3$ to be between 2 to $\infty$, such that we obtain $N_c = 2.15 \pm 0.17$, representing controllable non-factorizable effects [26] with $(a_1, a_2) = (1.00 \pm 0.01, 0.18 \pm 0.05)$ from Eq. (8). We list the branching
ratios of all possible two-body anti-triplet \( b \)-baryon decays in Table II and Table III, where the uncertainties are fitted with those from \((\rho, \eta, N_c)\), the decay constants and \(|\sqrt{3/2C_{||}}|\).

The decay branching ratios in Table III are given by \( a_1 \) with \( N_c = (2.15 \pm 0.17, \infty) \) as the theoretical inputs to demonstrate the insensitive non-factorizable effects. Note that \( N_c = 2.15 \pm 0.17 \) is fitted from \( B(\Lambda_b \rightarrow J/\psi \Lambda) \) and \( B(\Xi^- \rightarrow J/\psi \Xi^-) \), while \( N_c = \infty \) results in \( a_1 \approx c_{1}^{eff} \), wildly used in the generalized factorization. As the first measurement for the color-allowed decay mode, the predicted \( B(\Lambda_b \rightarrow D_s^- p) = (1.8 \pm 0.3) \times 10^{-5} \) or \( (2.5 \pm 0.4) \times 10^{-5} \) in Table III seems to disagree with the data in Eq. (9). Nonetheless, the predicted numbers driven by \( a_1 \) can be reliable as it is insensitive to the non-factorizable effects, whereas the data with the upper bound has a large uncertainty. Despite the color-allowed modes, the decay branching ratios of \( D^{(*)-} \) are found to be 30 times smaller than the \( D_s^{(*)-} \) counterparts. This can be simply understood by the relation of \((V_{cd}/V_{cs})^2(f_{D^{(*)}}/f_{D_s^{(*)}})^2 \simeq 0.03\). It is also interesting to note that the vector meson modes are 2 times as large as their pseudoscalar meson counterparts.

For the decay modes driven by \( a_2 \) as shown in Table III, we only list the results with \( a_2 = 0.18 \pm 0.05 \) (\( N_c = 2.15 \pm 0.17 \)). The reason is that \( a_2 \approx c_{2}^{eff} = -0.365 \) with \( N_c = \infty \) yields \( B(\Lambda_b \rightarrow J/\psi \Lambda) = (1.4 \pm 0.2) \times 10^{-3} \) and \( B(\Xi^- \rightarrow J/\psi \Xi^-) = (2.1 \pm 0.3) \times 10^{-3} \), which are in disagreement with the data in Eq. (9), demonstrating that the decays are sensitive

### Table II

| \( M_c = \) | \( D^- \) | \( D^{*-} \) |
|---|---|---|
| \( B(\Lambda_b \rightarrow pM_c) \) | (6.0 \( \pm \) 1.0, 8.2 \( \pm \) 1.4) \( \times \) 10\(^{-7}\) | (1.2 \( \pm \) 0.3, 1.6 \( \pm \) 0.4) \( \times \) 10\(^{-6}\) |
| \( B(\Xi^-_b \rightarrow \Lambda M_c) \) | (1.1 \( \pm \) 0.2, 1.5 \( \pm \) 0.2) \( \times \) 10\(^{-7}\) | (2.2 \( \pm \) 0.6, 3.0 \( \pm \) 0.8) \( \times \) 10\(^{-7}\) |
| \( B(\Xi^-_b \rightarrow \Sigma^0 M_c) \) | (3.3 \( \pm \) 0.5, 4.5 \( \pm \) 0.7) \( \times \) 10\(^{-7}\) | (6.6 \( \pm \) 1.6, 9.0 \( \pm \) 2.2) \( \times \) 10\(^{-7}\) |
| \( B(\Xi^-_b \rightarrow \Sigma^+ M_c) \) | (6.3 \( \pm \) 1.0, 8.6 \( \pm \) 1.4) \( \times \) 10\(^{-7}\) | (1.3 \( \pm \) 0.3, 1.7 \( \pm \) 0.4) \( \times \) 10\(^{-6}\) |

| \( M_c = \) | \( D_s^- \) | \( D_s^{*-} \) |
|---|---|---|
| \( B(\Lambda_b \rightarrow pM_c) \) | (1.8 \( \pm \) 0.3, 2.5 \( \pm \) 0.4) \( \times \) 10\(^{-5}\) | (3.5 \( \pm \) 0.9, 4.7 \( \pm \) 1.2) \( \times \) 10\(^{-5}\) |
| \( B(\Xi^-_b \rightarrow \Lambda M_c) \) | (3.4 \( \pm \) 0.5, 4.6 \( \pm \) 0.7) \( \times \) 10\(^{-6}\) | (6.4 \( \pm \) 1.6, 8.8 \( \pm \) 2.2) \( \times \) 10\(^{-6}\) |
| \( B(\Xi^-_b \rightarrow \Sigma^0 M_c) \) | (9.9 \( \pm \) 1.5, 13.6 \( \pm \) 2.1) \( \times \) 10\(^{-6}\) | (1.9 \( \pm \) 0.5, 2.6 \( \pm \) 0.6) \( \times \) 10\(^{-5}\) |
| \( B(\Xi^-_b \rightarrow \Sigma^+ M_c) \) | (1.9 \( \pm \) 0.3, 2.6 \( \pm \) 0.4) \( \times \) 10\(^{-5}\) | (3.6 \( \pm \) 0.9, 4.9 \( \pm \) 1.2) \( \times \) 10\(^{-5}\) |
to the non-factorizable effects. From Table III we see that both $\mathcal{B}(A_b \to J/\psi \Lambda)$ and $\mathcal{B}(\Xi_b^- \to J/\psi \Xi^-)$ are reproduced to agree with the data in Eq. (3) within errors. Note that $\mathcal{B}(A_b \to J/\psi \Lambda)/\mathcal{B}(\Xi_b^- \to J/\psi \Xi^-) \simeq 0.65$ in our calculation results from $(C|\|^2/(\sqrt{3/2}C|\|^2) \simeq 0.67$ as the ratio of their form factors in Table I which is in accordance with the $SU(3)$ flavor and $SU(2)$ spin symmetries. The more precise measurement of this ratio in the future will test the validity of the symmetries. As $\mathcal{B}(\Xi^- \to J/\psi \Sigma^-)$ is fitted by $C_3 = 2$ as shown in the right-bottom of Table II. In contrast, the neutral $D^{(*)0}$ modes fitted by $N_c = 2.15 \pm 0.17$.

| $M_c$ = | $D^0$ | $D^{*0}$ |
|---------|--------|--------|
| $\mathcal{B}(\Xi_b^- \to \Sigma^- M_c)$ | $(5.3 \pm 3.3) \times 10^{-5}$ | $(1.1 \pm 0.7) \times 10^{-4}$ |
| $\mathcal{B}(\Xi_b^0 \to \Lambda^0 M_c)$ | $(8.6 \pm 5.3) \times 10^{-6}$ | $(1.7 \pm 1.1) \times 10^{-5}$ |
| $\mathcal{B}(\Xi_b^0 \to \Sigma^0 M_c)$ | $(2.5 \pm 1.6) \times 10^{-5}$ | $(5.0 \pm 3.4) \times 10^{-5}$ |
| $\mathcal{B}(A_b \to \Lambda M_c)$ | $(1.6 \pm 1.0) \times 10^{-6}$ | $(3.3 \pm 2.2) \times 10^{-6}$ |
| $\mathcal{B}(A_b \to \Sigma M_c)$ | 0 | 0 |
| $\mathcal{B}(\Xi_b^- \to \Xi^- M_c)$ | $(2.7 \pm 1.7) \times 10^{-6}$ | $(5.5 \pm 3.6) \times 10^{-6}$ |
| $\mathcal{B}(\Xi_b^0 \to \Xi^0 M_c)$ | $(2.6 \pm 1.6) \times 10^{-6}$ | $(5.2 \pm 3.5) \times 10^{-6}$ |

| $M_c$ = | $\eta_c$ | $J/\psi$ |
|---------|--------|--------|
| $\mathcal{B}(\Xi_b^- \to \Sigma^- M_c)$ | $(1.4 \pm 0.8) \times 10^{-5}$ | $(2.9 \pm 1.8) \times 10^{-5}$ |
| $\mathcal{B}(\Xi_b^0 \to \Lambda^0 M_c)$ | $(2.3 \pm 1.4) \times 10^{-6}$ | $(4.7 \pm 2.9) \times 10^{-6}$ |
| $\mathcal{B}(\Xi_b^0 \to \Sigma^0 M_c)$ | $(6.6 \pm 4.1) \times 10^{-6}$ | $(1.4 \pm 0.8) \times 10^{-5}$ |
| $\mathcal{B}(A_b \to \Lambda M_c)$ | $(1.5 \pm 0.9) \times 10^{-4}$ | $(3.3 \pm 2.0) \times 10^{-4}$ |
| $\mathcal{B}(A_b \to \Sigma M_c)$ | 0 | 0 |
| $\mathcal{B}(\Xi_b^- \to \Xi^- M_c)$ | $(2.4 \pm 1.5) \times 10^{-4}$ | $(5.1 \pm 3.2) \times 10^{-4}$ |
| $\mathcal{B}(\Xi_b^0 \to \Xi^0 M_c)$ | $(2.3 \pm 1.4) \times 10^{-4}$ | $(4.9 \pm 3.0) \times 10^{-4}$ |
via the $b \to c\bar{d}$ transition have the branching ratios of order $10^{-6}$ caused by the suppression of $(V_{cb}V_{cd})^2/(V_{cb}V_{cs})^2 \simeq 0.225^2$. Finally, we remark that $\mathcal{B}(\Xi_b^+ \to \Xi^- M_c) \simeq \mathcal{B}(\Xi_b^0 \to \Xi^0 M_c)$ is due to the isospin symmetry.

### IV. CONCLUSIONS

In sum, we have studied all possible anti-triplet $B_b$ decays of the two-body charmful $B_b \to B_n M_c$ decays. We have found $\mathcal{B}(\Lambda_b \to D^- s) = (1.8 \pm 0.3) \times 10^{-5}$, which is within the measured upper bound of $\mathcal{B}(\Lambda_b \to D^- s) < 4.8(5.3) \times 10^{-4}$ at 90% (95%) C.L., and reproduced $\mathcal{B}(\Lambda_b \to J/\psi \Lambda) = (3.3 \pm 2.0) \times 10^{-4}$ and $\mathcal{B}(\Xi_b^- \to J/\psi \Xi^-) = (5.1 \pm 3.2) \times 10^{-4}$ in agreement with the data. Moreover, we have predicted $\mathcal{B}(\Lambda_b \to \Lambda \eta_c) = (1.5 \pm 0.9) \times 10^{-4}$, $\mathcal{B}(\Xi_b^- \to \Xi^- \eta_c) = (2.4 \pm 1.5) \times 10^{-4}$, and $\mathcal{B}(\Xi_b^0 \to \Xi^0 \eta_c, \Xi^0 J/\psi) = (2.3 \pm 1.4, 4.9 \pm 3.0) \times 10^{-4}$, which are accessible to the experiments at the LHCb, while $\mathcal{B}(\Xi_b^- \to \Xi^- M_c) \simeq \mathcal{B}(\Xi_b^0 \to \Xi^0 M_c)$ is due to the isospin symmetry.

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