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Charmful two-body antitriplet b-baryon decays

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We study the charmful decays of the two-body $\mathcal{B}_b \to \mathcal{B}_n M_c$ decays, where $\mathcal{B}_b$ represents the antitriplet of $(\Lambda_b, \Xi_b^0, \Xi_b^-)$, $\mathcal{B}_n$ stands for the baryon octet, and M_c denotes the charmed meson of $D_{(s)}^{(*)}$, η_c and J/ψ. 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 $\mathcal{B}(\Xi_b^0 \to \Xi^0\eta_c, \Xi^0 J/\psi) = (2.3 \pm 1.4, 4.9 \pm 3.0) \times 10^{-4}$, which are accessible to the experiments at the LHCb.

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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 through the $\Lambda_b \to \Lambda_c$ transition along with the recoiled mesons K^-, π^-, D_s^-, and D^-, respectively, such that one may use the factorization ansatz to get the fractions of the branching ratios as

$$\mathcal{R}_{K/\pi} \equiv \frac{\mathcal{B}(\Lambda_b \to \Lambda_c^+ K^-)}{\mathcal{B}(\Lambda_b \to \Lambda_c^+ \pi^-)} \simeq \frac{(|V_{us}| f_K)^2}{(|V_{ud}| f_\pi)^2} = 0.073,$$
$$\mathcal{R}_{D/D_s} \equiv \frac{\mathcal{B}(\Lambda_b \to \Lambda_c^+ D^-)}{\mathcal{B}(\Lambda_b \to \Lambda_c^+ D_s^-)} \simeq \frac{(|V_{cd}| f_D)^2}{(|V_{cs}| f_{D_s})^2} = 0.034, \quad (1)$$

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

$$\mathcal{R}_{K/\pi} = 0.0731 \pm 0.0016 \pm 0.0016,$$
$$\mathcal{R}_{D/D_s} = 0.042 \pm 0.003 \pm 0.003. \quad (2)$$

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

$$\mathcal{B}(\Lambda_b \to D_s^- p) = (2.7 \pm 1.4 \pm 0.2 \pm 0.7 \pm 0.1 \pm 0.1) \times 10^{-4} \quad \text{or}$$
$$< 4.8(5.3) \times 10^{-4} \quad \text{at } 90\%\ (95\%)\ \text{C.L.},$$
$$\mathcal{B}(\Lambda_b \to J/\psi\Lambda) = (3.0 \pm 1.1) \times 10^{-4},$$
$$\mathcal{B}(\Xi_b^- \to J/\psi\Xi^-) = (2.0 \pm 0.9) \times 10^{-4}, \quad (3)$$

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 $\mathcal{B}(\Xi_b^- \to J/\psi\Xi^-) f_{\Xi_b^-} = (1.02^{+0.26}_{-0.21}) \times 10^{-5}$, 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 Λ_b and Ξ_b [7], respectively. Nonetheless, for these $\mathcal{B}_b \to \mathcal{B}_n M_c$ 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 $\mathcal{B}_b \to \mathcal{B}_n M_c$, in this article we systematically analyze the branching ratios for all possible $\mathcal{B}_b \to \mathcal{B}_n M_c$ decays, and compare them with the experimental data at the B-factories, as well as the LHCb, where $\mathcal{B}_b$, $\mathcal{B}_n$, and M_c correspond to the antitriplet 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] are based on the factorization approach, the amplitudes for the two-body charmful b-baryon decays are presented in terms of the decaying process of the $\mathcal{B}_b \to \mathcal{B}_n$ transition with the recoiled charmed meson M_c. From Fig. 1(a), the amplitudes of $\mathcal{B}_b \to \mathcal{B}_n M_c$ via the quark-level $b \to u\bar{c}q$ transition are factorized as

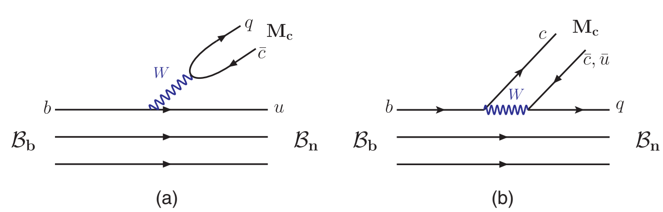

FIG. 1 (color online). Diagrams for two-body charmful $\mathcal{B}_b \to \mathcal{B}_n M_c$ decays.

$$\mathcal{A}_1(\mathcal{B}_b \to \mathcal{B}_n M_c) = \frac{G_F}{\sqrt{2}} V_{ub} V_{cq}^* a_1 \langle M_c|\bar{q}\gamma^\mu(1-\gamma_5)c|0\rangle \times \langle \mathcal{B}_n|\bar{u}\gamma_\mu(1-\gamma_5)b|\mathcal{B}_b\rangle, \quad (4)$$

where G_F is the Fermi constant, $V_{ub,cq}$ are the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements, and the explicit decay modes are

$$\Lambda_b \to pM_c, \qquad \Xi_b^- \to \Lambda(\Sigma^0)M_c, \qquad \Xi_b^0 \to \Sigma^+ M_c \quad (5)$$

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

$$\mathcal{A}_2(\mathcal{B}_b \to \mathcal{B}_n M_c) = \frac{G_F}{\sqrt{2}} V_{cb} V_{q_1 q}^* a_2 \langle M_c|\bar{c}\gamma^\mu(1-\gamma_5)q_1|0\rangle \times \langle \mathcal{B}_n|\bar{q}\gamma_\mu(1-\gamma_5)b|\mathcal{B}_b\rangle, \quad (6)$$

with $q_1 = u$ for $M_c = D^{(*)0}$ and $q_1 = c$ for $M_c = \eta_c$ and J/ψ, where the decays of $\mathcal{B}_b \to \mathcal{B}_n M_c$ are

$$\begin{aligned} &\Lambda_b \to nM_c, \qquad \Xi_b^- \to \Sigma^- M_c, \\ &\Xi_b^0 \to \Lambda(\Sigma^0)M_c \quad \text{for } q_2 = d, \\ &\Lambda_b \to \Lambda(\Sigma^0)M_c, \qquad \Xi_b^- \to \Xi^- M_c, \\ &\Xi_b^0 \to \Xi^0 M_c \quad \text{for } q_2 = s. \end{aligned} \quad (7)$$

In this article, we exclude the study of $\Lambda_b \to nM_c$ due to the elusive neutron in the B-factories. The amplitudes $\mathcal{A}_{1,2}$ 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]

$$a_1 = c_1^{\text{eff}} + \frac{c_2^{\text{eff}}}{N_c}, \qquad a_2 = c_2^{\text{eff}} + \frac{c_1^{\text{eff}}}{N_c}, \quad (8)$$

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 $2 \to \infty$ to account for the nonfactorizable 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

$$\begin{aligned} \langle P_c|A_\mu^c|0\rangle &= -if_{P_c} q_\mu, \\ \langle V_c|V_\mu^c|0\rangle &= m_{V_c} f_{V_c} \varepsilon_\mu^*, \end{aligned} \quad (9)$$

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_μ and ε_μ^* are the four-momentum and polarization, respectively. Those of the $\mathcal{B}_b \to \mathcal{B}_n$ baryon transition in Eq. (4) have the general forms

$$\begin{aligned} \langle \mathcal{B}_n|\bar{q}\gamma_\mu b|\mathcal{B}_b\rangle &= \bar{u}_{\mathcal{B}_n}\left[f_1\gamma_\mu + \frac{f_2}{m_{\mathcal{B}_b}} i\sigma_{\mu\nu}q^\nu + \frac{f_3}{m_{\mathcal{B}_b}} q_\mu\right] u_{\mathcal{B}_b}, \\ \langle \mathcal{B}_n|\bar{q}\gamma_\mu\gamma_5 b|\mathcal{B}_b\rangle &= \bar{u}_{\mathcal{B}_n}\left[g_1\gamma_\mu + \frac{g_2}{m_{\mathcal{B}_b}} i\sigma_{\mu\nu}q^\nu + \frac{g_3}{m_{\mathcal{B}_b}} q_\mu\right]\gamma_5 u_{\mathcal{B}_b}, \end{aligned} \quad (10)$$

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

$$\langle \mathcal{B}_{n,\uparrow+\downarrow}|J_{\mu,R}^q|\mathcal{B}_{b,\uparrow+\downarrow}\rangle = \bar{u}_{\mathcal{B}_n}\left[\gamma_\mu\frac{1+\gamma_5}{2}G^\uparrow(q^2) + \gamma_\mu\frac{1-\gamma_5}{2}G^\downarrow(q^2)\right]u_{\mathcal{B}_b}, \quad (11)$$

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

$$\begin{aligned} G^\uparrow(q^2) &= e_\parallel^\uparrow G_\parallel(q^2) + e_{\overline{\parallel}}^\uparrow G_{\overline{\parallel}}(q^2), \\ G^\downarrow(q^2) &= e_\parallel^\downarrow G_\parallel(q^2) + e_{\overline{\parallel}}^\downarrow G_{\overline{\parallel}}(q^2), \end{aligned} \quad (12)$$

with the constants $e_{\parallel(\overline{\parallel})}^\uparrow$ and $e_{\parallel(\overline{\parallel})}^\downarrow$ to sum over the chiral charges via the $\mathcal{B}_b \to \mathcal{B}_n$ transition, given by

$$\begin{aligned} e_\parallel^\uparrow &= \langle \mathcal{B}_{n,\uparrow}|\mathbf{Q}_\parallel|\mathcal{B}_{b,\uparrow}\rangle, \qquad e_{\overline{\parallel}}^\uparrow = \langle \mathcal{B}_{n,\uparrow}|\mathbf{Q}_{\overline{\parallel}}|\mathcal{B}_{b,\uparrow}\rangle, \\ e_\parallel^\downarrow &= \langle \mathcal{B}_{n,\downarrow}|\mathbf{Q}_\parallel|\mathcal{B}_{b,\downarrow}\rangle, \qquad e_{\overline{\parallel}}^\downarrow = \langle \mathcal{B}_{n,\downarrow}|\mathbf{Q}_{\overline{\parallel}}|\mathcal{B}_{b,\downarrow}\rangle. \end{aligned} \quad (13)$$

Note that $\mathbf{Q}_{\parallel(\overline{\parallel})} = \sum_i Q_{\parallel(\overline{\parallel})}(i)$ with $i = 1, 2, 3$ as the the chiral charge operators are from $Q_R^q \equiv J_{0,R}^q = q_R^\dagger b_R$, converting the b quark in $|\mathcal{B}_{b,\uparrow,\downarrow}\rangle$ into the q one, while the converted q quark can be parallel or antiparallel to the $\mathcal{B}_b$'s helicity, denoted as the subscript $\parallel$ or $\overline{\parallel}$. By comparing Eq. (10) with Eqs. (11)–(13), we obtain

$$\begin{aligned} f_1 &= (e_\parallel^\uparrow + e_\parallel^\downarrow)G_\parallel + (e_{\overline{\parallel}}^\uparrow + e_{\overline{\parallel}}^\downarrow)G_{\overline{\parallel}}, \\ g_1 &= (e_\parallel^\uparrow - e_\parallel^\downarrow)G_\parallel + (e_{\overline{\parallel}}^\uparrow - e_{\overline{\parallel}}^\downarrow)G_{\overline{\parallel}}, \end{aligned} \quad (14)$$

with $f_{2,3} = 0$ and $g_{2,3} = 0$ due to the helicity conservation, as those derived in Refs. [8,10,23]. It is interesting to see that, like the helicity-flip terms, the theoretical calculations

TABLE I. Relations between the transition matrix elements.

$\langle\mathcal{B}_n\vert(\bar{q}b)\vert\mathcal{B}_b\rangle$	$f_1(0)=g_1(0)$
$\langle p\vert(\bar{u}b)\vert\Lambda_b\rangle$	$-\sqrt{\frac{3}{2}}C_\parallel$
$\langle\Lambda\vert(\bar{u}b)\vert\Xi_b^-\rangle$	$\frac{1}{2}C_\parallel$
$\langle\Sigma^0\vert(\bar{u}b)\vert\Xi_b^-\rangle$	$-\sqrt{\frac{3}{4}}C_\parallel$
$\langle\Sigma^+\vert(\bar{u}b)\vert\Xi_b^0\rangle$	$-\sqrt{\frac{3}{2}}C_\parallel$
$\langle\Sigma^-\vert(\bar{d}b)\vert\Xi_b^-\rangle$	$\sqrt{\frac{3}{2}}C_\parallel$
$\langle\Lambda\vert(\bar{d}b)\vert\Xi_b^0\rangle$	$-\frac{1}{2}C_\parallel$
$\langle\Sigma^0\vert(\bar{d}b)\vert\Xi_b^0\rangle$	$\sqrt{\frac{3}{4}}C_\parallel$
$\langle\Lambda\vert(\bar{s}b)\vert\Lambda_b\rangle$	$C_\parallel$
$\langle\Sigma^0\vert(\bar{s}b)\vert\Lambda_b\rangle$	0
$\langle\Xi^-\vert(\bar{s}b)\vert\Xi_b^-\rangle$	$\sqrt{\frac{3}{2}}C_\parallel$
$\langle\Xi^0\vert(\bar{s}b)\vert\Xi_b^0\rangle$	$-\sqrt{\frac{3}{2}}C_\parallel$

from the loop contributions to $f_{2,3}$ $(g_{2,3})$ indeed result in the values being 1 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 [3]

$$f_1(q^2)=\frac{f_1(0)}{(1-q^2/m_{\mathcal{B}_b}^2)^2},\qquad g_1(q^2)=\frac{g_1(0)}{(1-q^2/m_{\mathcal{B}_b}^2)^2},\tag{15}$$

such that it is reasonable to parametrize the chiral form factors to be $(1-q^2/m_{\mathcal{B}_b}^2)^2 G_{\parallel(\bar{\parallel})}=C_{\parallel(\bar{\parallel})}$. Subsequently, from

$$\begin{aligned}&\left(e_\parallel^\uparrow,e_\parallel^\downarrow,e_{\bar\parallel}^\uparrow,e_{\bar\parallel}^\downarrow\right)=(-\sqrt{3/2},0,0,0)\quad\text{for }\langle p|J_{\mu,R}^u|\Lambda_b\rangle,\\&\left(e_\parallel^\uparrow,e_\parallel^\downarrow,e_{\bar\parallel}^\uparrow,e_{\bar\parallel}^\downarrow\right)=(1,0,0,0)\quad\text{for }\langle\Lambda|J_{\mu,R}^s|\Lambda_b\rangle,\\&\left(e_\parallel^\uparrow,e_\parallel^\downarrow,e_{\bar\parallel}^\uparrow,e_{\bar\parallel}^\downarrow\right)=(0,0,0,0)\quad\text{for }\langle\Sigma^0|J_{\mu,R}^s|\Lambda_b\rangle,\end{aligned}\tag{16}$$

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 antitriplet b baryons, $(\Xi_b^-,\Xi_b^0,\Lambda_b^0)$ shown in Table I, we find that the relation of $f_1=g_1$ is uniquely determined for the antitriplet b-baryon transitions.

III. NUMERICAL RESULTS AND DISCUSSIONS

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

$$\begin{aligned}&(V_{ub},V_{cb})=(A\lambda^3(\rho-i\eta),A\lambda^2),\\&(V_{cd}=-V_{us},V_{cs}=V_{ud})=(-\lambda,1-\lambda^2/2),\end{aligned}\tag{17}$$

with $(\lambda,A,\rho,\eta)=(0.225,0.814,0.120\pm0.022,0.362\pm0.013)$. The meson decay constants are adopted as $(f_{\eta_c},f_{J/\psi})=(387\pm7,418\pm9)$ MeV [24], $(f_D,f_{D_s})=(204.6\pm5.0,257.5\pm4.6)$ MeV [4], and $(f_{D^*},f_{D_s^*})=(252.2\pm22.7,305.5\pm27.3)$ 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 $|\sqrt{3/2}C_\parallel|=0.136\pm0.009$ for $\langle p|\bar{u}\gamma_\mu(\gamma_5)b|\Lambda_b\rangle$, which is consistent with the value of 0.14 ± 0.03 in the light-cone sum rules [23] and the theoretical calculations in Refs. [8,10]. With $\mathcal{B}(\Lambda_b\to J/\psi\Lambda)$ and $\mathcal{B}(\Xi_b^-\to J/\psi\Xi^-)$ in Eq. (3) as the experimental inputs, we can estimate the nonfactorizable effects by deviating the color number $N_c=3$ to be between 2 and ∞, such that we obtain $N_c=2.15\pm0.17$, representing controllable nonfactorizable effects [26] with $(a_1,a_2)=(1.00\pm0.01,0.18\pm0.05)$ from Eq. (8). We list the branching ratios of all possible two-body antitriplet b-baryon decays in Tables II–III, where the uncertainties are fitted with those from (ρ,η,N_c), the decay constants and $|\sqrt{3/2}C_\parallel|$.

The decay branching ratios in Table II are given by a_1 with $N_c=(2.15\pm0.17,\infty)$ as the theoretical inputs to demonstrate the insensitive nonfactorizable effects. Note that $N_c=2.15\pm0.17$ is fitted from $\mathcal{B}(\Lambda_b\to J/\psi\Lambda)$ and $\mathcal{B}(\Xi_b^-\to J/\psi\Xi^-)$, while $N_c=\infty$ results in $a_1\simeq c_1^{\rm eff}$, wildly used in the generalized factorization. As the first measurement for the color-allowed decay mode, the predicted $\mathcal{B}(\Lambda_b\to D_s^-p)=(1.8\pm0.3)\times10^{-5}$ or $(2.5\pm0.4)\times10^{-5}$ in Table II seems to disagree with the data in Eq. (3). Nonetheless, the predicted numbers driven by a_1 can be reliable as it is insensitive to the nonfactorizble effects, whereas the data with the upper bound have 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 $(V_{cd}/V_{cs})^2(f_{D^{(*)}}/f_{D_s^{(*)}})^2\simeq0.03$. It is also interesting to note that the vector meson modes are two 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 $a_2=0.18\pm0.05$ $(N_c=2.15\pm0.17)$. The reason is that $a_2\simeq c_2^{\rm eff}=-0.365$ with $N_c=\infty$ yields $\mathcal{B}(\Lambda_b\to J/\psi\Lambda)=(1.4\pm0.2)\times10^{-3}$ and $\mathcal{B}(\Xi_b^-\to J/\psi\Xi^-)=(2.1\pm0.3)\times10^{-3}$, which are in disagreement with the data in Eq. (3), demonstrating that the decays are

TABLE II. The branching ratios of all possible two-body antitriplet b-baryon decays with a_1 fitted by $N_c = (2.15 \pm 0.17, \infty)$.

$M_c =$	D^-	D^{*-}
$\mathcal{B}(\Lambda_b \to pM_c)$	$(6.0 \pm 1.0, 8.2 \pm 1.4) \times 10^{-7}$	$(1.2 \pm 0.3, 1.6 \pm 0.4) \times 10^{-6}$
$\mathcal{B}(\Xi_b^- \to \Lambda M_c)$	$(1.1 \pm 0.2, 1.5 \pm 0.2) \times 10^{-7}$	$(2.2 \pm 0.6, 3.0 \pm 0.8) \times 10^{-7}$
$\mathcal{B}(\Xi_b^- \to \Sigma^0 M_c)$	$(3.3 \pm 0.5, 4.5 \pm 0.7) \times 10^{-7}$	$(6.6 \pm 1.6, 9.0 \pm 2.2) \times 10^{-7}$
$\mathcal{B}(\Xi_b^0 \to \Sigma^+ M_c)$	$(6.3 \pm 1.0, 8.6 \pm 1.4) \times 10^{-7}$	$(1.3 \pm 0.3, 1.7 \pm 0.4) \times 10^{-6}$
$M_c =$	D_s^-	D_s^{*-}
$\mathcal{B}(\Lambda_b \to pM_c)$	$(1.8 \pm 0.3, 2.5 \pm 0.4) \times 10^{-5}$	$(3.5 \pm 0.9, 4.7 \pm 1.2) \times 10^{-5}$
$\mathcal{B}(\Xi_b^- \to \Lambda M_c)$	$(3.4 \pm 0.5, 4.6 \pm 0.7) \times 10^{-6}$	$(6.4 \pm 1.6, 8.8 \pm 2.2) \times 10^{-6}$
$\mathcal{B}(\Xi_b^- \to \Sigma^0 M_c)$	$(9.9 \pm 1.5, 13.6 \pm 2.1) \times 10^{-6}$	$(1.9 \pm 0.5, 2.6 \pm 0.6) \times 10^{-5}$
$\mathcal{B}(\Xi_b^0 \to \Sigma^+ M_c)$	$(1.9 \pm 0.3, 2.6 \pm 0.4) \times 10^{-5}$	$(3.6 \pm 0.9, 4.9 \pm 1.2) \times 10^{-5}$

sensitive to the nonfactorizable effects. From Table III, we see that both $\mathcal{B}(\Lambda_b \to J/\psi\Lambda)$ and $\mathcal{B}(\Xi_b^- \to J/\psi\Xi^-)$ are reproduced to agree with the data in Eq. (3) within errors. Note that $\mathcal{B}(\Lambda_b \to J/\psi\Lambda)/\mathcal{B}(\Xi_b^- \to J/\psi\Xi^-) \simeq 0.65$ in our calculation results from $(C_\parallel)^2/(\sqrt{3/2}C_\parallel)^2 \simeq 0.67$ 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 $\mathcal{B}(\Xi_b^- \to J/\psi\Xi^-) = O(10^{-4})$, we emphasize that more experimental searches should be done for the two-body Ξ_b decays, while most of the recent observations are from the Λ_b decays. Since the result of $\mathcal{B}(\Lambda_b \to \Sigma^0 M_c) = 0$ is from the $SU(3)$ flavor and $SU(2)$ spin symmetries as well as the heavy flavor symmetry, a non-zero measurement will break the symmetries. Through the $b \to c\bar{c}s$ transition at the quark level, $\mathcal{B}(\Lambda_b \to \Lambda M_c)$, $\mathcal{B}(\Xi_b^- \to \Xi^- M_c)$, and $\mathcal{B}(\Xi_b^0 \to \Xi^0 M_c)$ with $M_c = \eta_c$ and J/ψ are all $O(10^{-4})$ as shown in the bottom right of Table II. In contrast, the neutral $D^{(*)0}$ modes via the $b \to c\bar{c}d$ 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.

TABLE III. The branching ratios of all possible two-body antitriplet b-baryon decays with a_2 fitted by $N_c = 2.15 \pm 0.17$.

$M_c =$	D^0	D^{*0}
$\mathcal{B}(\Xi_b^- \to \Sigma^- M_c)$	$(5.3 \pm 3.3) \times 10^{-5}$	$(1.1 \pm 0.7) \times 10^{-4}$
$\mathcal{B}(\Xi_b^0 \to \Lambda^0 M_c)$	$(8.6 \pm 5.3) \times 10^{-6}$	$(1.7 \pm 1.1) \times 10^{-5}$
$\mathcal{B}(\Xi_b^0 \to \Sigma^0 M_c)$	$(2.5 \pm 1.6) \times 10^{-5}$	$(5.0 \pm 3.4) \times 10^{-5}$
$\mathcal{B}(\Lambda_b \to \Lambda M_c)$	$(1.6 \pm 1.0) \times 10^{-6}$	$(3.3 \pm 2.2) \times 10^{-6}$
$\mathcal{B}(\Lambda_b \to \Sigma^0 M_c)$	0	0
$\mathcal{B}(\Xi_b^- \to \Xi^- M_c)$	$(2.7 \pm 1.7) \times 10^{-6}$	$(5.5 \pm 3.6) \times 10^{-6}$
$\mathcal{B}(\Xi_b^0 \to \Xi^0 M_c)$	$(2.6 \pm 1.6) \times 10^{-6}$	$(5.2 \pm 3.5) \times 10^{-6}$
$M_c =$	η_c	J/ψ
$\mathcal{B}(\Xi_b^- \to \Sigma^- M_c)$	$(1.4 \pm 0.8) \times 10^{-5}$	$(2.9 \pm 1.8) \times 10^{-5}$
$\mathcal{B}(\Xi_b^0 \to \Lambda^0 M_c)$	$(2.3 \pm 1.4) \times 10^{-6}$	$(4.7 \pm 2.9) \times 10^{-6}$
$\mathcal{B}(\Xi_b^0 \to \Sigma^0 M_c)$	$(6.6 \pm 4.1) \times 10^{-6}$	$(1.4 \pm 0.8) \times 10^{-5}$
$\mathcal{B}(\Lambda_b \to \Lambda M_c)$	$(1.5 \pm 0.9) \times 10^{-4}$	$(3.3 \pm 2.0) \times 10^{-4}$
$\mathcal{B}(\Lambda_b \to \Sigma^0 M_c)$	0	0
$\mathcal{B}(\Xi_b^- \to \Xi^- M_c)$	$(2.4 \pm 1.5) \times 10^{-4}$	$(5.1 \pm 3.2) \times 10^{-4}$
$\mathcal{B}(\Xi_b^0 \to \Xi^0 M_c)$	$(2.3 \pm 1.4) \times 10^{-4}$	$(4.9 \pm 3.0) \times 10^{-4}$

IV. CONCLUSIONS

In sum, we have studied all possible antitriplet $\mathcal{B}_b$ decays of the two-body charmful $\mathcal{B}_b \to \mathcal{B}_n M_c$ decays. We have found $\mathcal{B}(\Lambda_b \to D_s^- p) = (1.8 \pm 0.3) \times 10^{-5}$, 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 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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