Antitriplet charmed baryon decays with SU(3) flavor symmetry

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(Received 11 January 2018; published 26 April 2018)

We study the decays of the antitriplet charmed baryon state $(\Xi_c^0, \Xi_c^+, \Lambda_c^+)$ based on the $SU(3)$ flavor symmetry. In particular, after predicting $\mathcal{B}(\Xi_c^0 \to \Xi^-\pi^+) = (15.7\pm0.7)\times10^{-3}$ and $\mathcal{B}(\Xi_c^+ \to \Xi^-\pi^+\pi^+) = (14.7\pm8.4)\times10^{-3}$, we extract that $\mathcal{B}(\Xi_c^0 \to \Lambda K^-\pi^+, \Lambda K^+K^-, \Xi^- e^+\nu_e) = (16.8\pm2.3, 0.45\pm0.11, 48.7\pm17.4)\times10^{-3}$ and $\mathcal{B}(\Xi_c^+ \to pK_s^0K_s^0, \Sigma^+K^-\pi^+, \Xi^0\pi^+\pi^0, \Xi^0 e^+\nu_e) = (1.3\pm0.8, 13.8\pm8.0, 33.8\pm21.9, 33.8^{+21.9}_{-22.6})\times10^{-3}$. We also find that $\mathcal{B}(\Xi_c^0 \to \Xi^0\eta, \Xi^0\eta') = (1.7^{+1.0}_{-1.7}, 8.6^{+11.0}_{-6.3})\times10^{-3}$, $\mathcal{B}(\Xi_c^0 \to \Lambda^0\eta, \Lambda^0\eta') = (1.6^{+1.2}_{-0.8}, 9.4^{+11.6}_{-6.8})\times10^{-4}$ and $\mathcal{B}(\Xi_c^+ \to \Sigma^+\eta, \Sigma^+\eta') = (28.4^{+8.2}_{-6.9}, 13.2^{+24.0}_{-11.9})\times10^{-4}$. These Ξ_c decays with the branching ratios of $O(10^{-4}\text{–}10^{-3})$ are clearly promising to be observed by the BESIII and LHCb experiments.

DOI: 10.1103/PhysRevD.97.073006

I. INTRODUCTION

In terms of the $SU(3)$ flavor $(SU(3)_f)$ symmetry, the Ξ_c decays should be in association with the Λ_c^+ ones as Ξ_c^0, Ξ_c^+, and Λ_c^+ are united as the lowest-lying antitriplet of the charmed baryon states $(\mathbf{B}_c)$. Nonetheless, in accordance with $f_{\Xi_c^+} + f_{\Xi_c^0} + f_{\Omega_c^0} \simeq 0.136 f_{\Lambda_c^+}$ estimated in Refs. [1,2], in which $f_{\mathbf{B}_c,\Omega_c^0}$ stand for the fragmentation fractions for the rates of the charmed baryon productions, the measurements of the Ξ_c decays are not easy tasks compared to the Λ_c^+ ones. For example, the two-body $\Lambda_c^+ \to \mathbf{B}_n M$ decays with $\mathbf{B}_n(M)$ the baryon (pseudoscalar meson) have been extensively studied by experiments. Interestingly, six decay Λ_c^+ decay modes have been recently reexamined or measured by BESIII [3,4]. In addition, LHCb has just observed the three-body $\Lambda_c^+ \to pMM$ decays [5], together with their CP-violating asymmetries [6]. However, not much progress has been made in the Ξ_c decays. In particular, none of the absolute branching fractions in the Ξ_c decays has been given yet. Instead, these decays are experimentally measured by relating the decays of $\Xi_c^+ \to \Xi^-\pi^+\pi^+$ or $\Xi_c^0 \to \Xi^-\pi^+$ and can only be determined once $f_{\Xi_c^{0,+}}$ [7] are known.

Since BESIII and LHCb are expected to search for all possible antitriplet charmed baryon decays, one can test whether or not the studies of $\Lambda_c^+ \to \mathbf{B}_n M$ can be applied to $\Xi_c^{0,+} \to \mathbf{B}_n M$. Theoretically, the factorization for the b baryon decays [8–13] does not work for the charmed baryon decays, which receive corrections by taking into account the nonfactorizable effects [14–19]. On the other hand, the possible b or c hadron decay modes can be examined by the $SU(3)_f$ symmetry [20–31]. Furthermore, the symmetry approach has been extended to explore the doubly and triply charmed baryon decays [31], which helps to establish the spectroscopies of the doubly and triply charmed baryon states [32], such as the to-be-confirmed Ξ_{cc}^+ state [33–38].

Moreover, to test the validity of the $SU(3)_f$ symmetry in the antitriplet charmed baryon decays, a complete numerical analysis for the decays is necessary. In fact, the decays of $\Lambda_c^+ \to \mathbf{B}_n M$ have been explained well by the global fit in Ref. [30], together with the predictions of $\mathcal{B}(\Xi_c^+ \to \Xi^0\pi^+) = (8.0\pm4.1)\times10^{-3}$ and $\mathcal{B}(\Xi_c^0 \to \Lambda^0\bar{K}^0) = (8.3\pm0.9)\times10^{-3}$, in agreement with the values of $(7.2\pm3.5, 8.3\pm3.7)\times10^{-3}$ extracted from the ratios of $\mathcal{B}(\Xi_c^+ \to \Xi^0\pi^+)/\mathcal{B}(\Xi_c^+ \to \Xi^0 e^+\nu_e)$ and $\mathcal{B}(\Xi_c^0 \to \Lambda^0\bar{K}^0)/\mathcal{B}(\Xi_c^0 \to \Xi^- e^+\nu_e)$, respectively [31].

In this report, we will systematically study the two-body weak $\Xi_c \to \mathbf{B}_n M$ decays based on the $SU(3)_f$ symmetry and give some specific numerical results, which can be tested in the future measurements by BESIII and LHCb. By taking the predicted $\mathcal{B}(\Xi_c^0 \to \Xi^-\pi^+)$ as the theoretical input, we will also estimate the branching ratios of other Ξ_c decays in the Particle Data Group (PDG) [7], which are related to $\Xi_c^0 \to \Xi^-\pi^+$.

II. FORMALISM

For the two-body antitriplet of the lowest-lying charmed baryon decays of $\mathbf{B}_c \to \mathbf{B}_n M$, where $\mathbf{B}_c = (\Xi_c^0, -\Xi_c^+, \Lambda_c^+)$ and $\mathbf{B}_n$ (M) are the baryon (pseudoscalar) octet states, the effective Hamiltonian responsible for the tree-level $c \to s u \bar{d}$, $c \to u q \bar{q}$, and $c \to d u \bar{s}$ transitions are given by [39]

$$\mathcal{H}_{\rm eff} = \sum_{i=+,-} \frac{G_F}{\sqrt{2}} c_i (V_{cs} V_{ud} O_i + V_{cd} V_{ud} O_i^{\dagger} + V_{cd} V_{us} O_i'), \tag{1}$$

with $q\bar{q} = d\bar{d}$ or $s\bar{s}$, G_F the Fermi constant, V_{ij} the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements, and $c_\pm$ the scale-dependent Wilson coefficients to take into account the subleading-order QCD corrections. The four-quark operators $O_\pm^{(\prime)}$ and $O_\pm^\dagger \equiv O_\pm^d - O_\pm^s$ in Eq. (1) can be written as

$$\begin{aligned} O_\pm &= \frac{1}{2}[(\bar{u}d)_{V-A}(\bar{s}c)_{V-A} \pm (\bar{s}d)_{V-A}(\bar{u}c)_{V-A}], \\ O_\pm^q &= \frac{1}{2}[(\bar{u}q)_{V-A}(\bar{q}c)_{V-A} \pm (\bar{q}q)_{V-A}(\bar{u}c)_{V-A}], \\ O'_\pm &= \frac{1}{2}[(\bar{u}s)_{V-A}(\bar{d}c)_{V-A} \pm (\bar{d}s)_{V-A}(\bar{u}c)_{V-A}], \end{aligned} \tag{2}$$

where $(\bar{q}_1 q_2)_{V-A} = \bar{q}_1 \gamma_\mu (1-\gamma_5) q_2$. By using $(V_{cs}V_{ud}, V_{cd}V_{ud}, V_{cd}V_{us}) \simeq (1, -s_c, -s_c^2)$ in Eq. (1) with $s_c \equiv \sin\theta_c = 0.2248$ [7] representing the well-known Cabbibo angle θ_c, the decays with $O_\pm$, $O_\pm^\dagger$, and $O'_\pm$ are the so-called Cabibbo-allowed, Cabibbo-suppressed, and doubly Cabibbo-suppressed processes, respectively. For instance, of the Cabibbo-allowed decay, $\mathcal{B}(\Lambda_c^+ \to p\bar{K}^0) = (3.16 \pm 0.16) \times 10^{-2}$ is measured to be 50 times larger than $\mathcal{B}(\Lambda_c^+ \to \Lambda K^+) = (6.1 \pm 1.2) \times 10^{-4}$, which is the Cabibbo-suppressed case, whereas none of the doubly Cabibbo-suppressed ones has been observed [7].

Without explicitly showing the Lorentz indices, the operators in Eq. (2) behave as $(\bar{q}^i q_k \bar{q}^j)c$, with $q_i = (u, d, s)$ as the triplet of 3, which can be decomposed as the irreducible forms under the $SU(3)_f$ symmetry, that is, $(\bar{3} \times 3 \times \bar{3})c = (\bar{3} + \bar{3}' + 6 + \overline{15})c$. Accordingly, (O_-, O_+) fall into the irreducible presentations of $(\mathcal{O}_6, \mathcal{O}_{\overline{15}})$, given by [25]

$$\mathcal{O}_6 = \frac{1}{2}(\bar{u}d\bar{s} - \bar{s}d\bar{u})c, \mathcal{O}_{\overline{15}} = \frac{1}{2}(\bar{u}d\bar{s} + \bar{s}d\bar{u})c, \tag{3}$$

which correspond to the tensor notations of $1/2\epsilon^{ijl}H(6)_{lk}$ and $H(\overline{15})_k^{ij}$, respectively, with (i, j, k) representing the quark indices and the nonzero entries being $H_{22}(6) = 2$ and $H_2^{13}(\overline{15}) = H_2^{31}(\overline{15}) = 1$. Note that $O_\pm^\dagger$ and $O'_\pm$ also have similar irreducible representations, resulting in the nonzero entries of $H_{23,32}(6) = -2s_c$, $H_2^{12,21}(\overline{15}) = -H_3^{13,31}(\overline{15}) = s_c$, $H_{33}(6) = 2s_c^2$, and $H_3^{12,21}(\overline{15}) = -s_c^2$ [25]. By using the bases of the $SU(3)_f$ symmetry, the effective Hamiltonian in Eq. (1) is transformed as

$$\mathcal{H}_{\rm eff} = \frac{G_F}{\sqrt{2}}\left[c_- \frac{\epsilon^{ijl}}{2} H(6)_{lk} + c_+ H(\overline{15})_k^{ij}\right], \tag{4}$$

where the individual nonzero entries of $H(6)_{lk}$ and $H(\overline{15})_k^{ij}$ that include $O_\mp$, $O_\mp^\dagger$, and $O'_\mp$ can be presented as the matrix forms:

$$H(6) = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 2 & -2s_c \\ 0 & -2s_c & 2s_c^2 \end{pmatrix},$$
$$H(\overline{15}) = \left(\begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \begin{pmatrix} 0 & s_c & 1 \\ s_c & 0 & 0 \\ 1 & 0 & 0 \end{pmatrix}, \begin{pmatrix} 0 & -s_c^2 & -s_c \\ -s_c^2 & 0 & 0 \\ -s_c & 0 & 0 \end{pmatrix}\right). \tag{5}$$

Correspondingly, the $\mathbf{B}_c$ antitriplet and $\mathbf{B}_n$ octet states are written as

$$\mathbf{B}_c = (\Xi_c^0, -\Xi_c^+, \Lambda_c^+),$$
$$\mathbf{B}_n = \begin{pmatrix} \frac{1}{\sqrt{6}}\Lambda + \frac{1}{\sqrt{2}}\Sigma^0 & \Sigma^+ & p \\ \Sigma^- & \frac{1}{\sqrt{6}}\Lambda - \frac{1}{\sqrt{2}}\Sigma^0 & n \\ \Xi^- & \Xi^0 & -\sqrt{\frac{2}{3}}\Lambda \end{pmatrix}. \tag{6}$$

The adding of the singlet η_1 to the octet (π, K, η_8) leads to the nonet of the pseudoscalar meson, given by [30],

$$M = \begin{pmatrix} \frac{1}{\sqrt{2}}(\pi^0 + c_\phi \eta + s_\phi \eta') & \pi^- & K^- \\ \pi^+ & \frac{-1}{\sqrt{2}}(\pi^0 - c_\phi \eta - s_\phi \eta') & \bar{K}^0 \\ K^+ & K^0 & -s_\phi \eta + c_\phi \eta' \end{pmatrix}, \tag{7}$$

where (η, η') are the mixtures of (η_1, η_8), with the mixing angle $\phi = (39.3 \pm 1.0)^\circ$ [40] for $(c_\phi, s_\phi) = (\cos\phi, \sin\phi)$.

The amplitudes of the $\mathbf{B}_c \to \mathbf{B}_n M$ decays via the effective Hamiltonian in Eq. (1) appear to be $\mathcal{A}(\mathbf{B}_c \to \mathbf{B}_n M) = \langle \mathbf{B}_n M | \mathcal{H}_{\rm eff} | \mathbf{B}_c \rangle$. Since $\mathcal{H}_{\rm eff}$, $\mathbf{B}_{c(n)}$, and M have been in the $SU(3)_f$ forms, the amplitudes of $\mathbf{B}_c \to \mathbf{B}_n M$ can be further derived as

$$\begin{aligned} \mathcal{A}(\mathbf{B}_c \to \mathbf{B}_n M) &= \langle \mathbf{B}_n M | \mathcal{H}_{\rm eff} | \mathbf{B}_c \rangle \\ &= \frac{G_F}{\sqrt{2}} T(\mathbf{B}_c \to \mathbf{B}_n M), \end{aligned} \tag{8}$$

with $T(\mathbf{B}_c \to \mathbf{B}_n M)$ given by [28]

$$
\begin{aligned}
T(\mathbf{B}_c \to \mathbf{B}_n M) &= T(\mathcal{O}_6) + T(\mathcal{O}_{\overline{15}}) \\
T(\mathcal{O}_6) &= a_1 H_{ij}(6) T^{ik} (\mathbf{B}_n)^l_k (M)^j_l \\
&\quad + a_2 H_{ij}(6) T^{ik} (M)^l_k (\mathbf{B}_n)^j_l \\
&\quad + a_3 H_{ij}(6) (\mathbf{B}_n)^i_k (M)^j_l T^{kl} \\
&\quad + h H_{ij}(6) T^{ik} (\mathbf{B}_n)^j_k (M)^l_l, \\
T(\mathcal{O}_{\overline{15}}) &= a_4 H^k_{li}(\overline{15}) (\mathbf{B}_c)^j (M)^i_j (\mathbf{B}_n)^l_k \\
&\quad + a_5 (\mathbf{B}_n)^i_j (M)^l_i H(\overline{15})^{jk}_l (\mathbf{B}_c)_k \\
&\quad + a_6 (\mathbf{B}_n)^k_l (M)^i_j H(\overline{15})^{jl}_i (\mathbf{B}_c)_k \\
&\quad + a_7 (\mathbf{B}_n)^l_i (M)^i_j H(\overline{15})^{jk}_l (\mathbf{B}_c)_k \\
&\quad + h' H^{jk}_i(\overline{15}) (\mathbf{B}_n)^i_k (M)^l_l (\mathbf{B}_c)_j, \qquad (9)
\end{aligned}
$$

where $T_{ij} \equiv (\mathbf{B}_c)_k \epsilon^{ijk}$; (c_-, c_+) have been absorbed into the $SU(3)$ parameters of (a_1, a_2, a_3, h) and (a_4, a_5, a_6, a_7, h'), respectively; and the $h^{(\prime)}$ terms correspond to the contributions from the singlet η_1. With the T amps expanded in Table I, we are enabled to relate all possible two-body $\mathbf{B}_c \to \mathbf{B}_n M$ decays with the $SU(3)_f$ parameters. To compute the branching ratios, we use the equation given by [7]

$$
\mathcal{B}(\mathbf{B}_c \to \mathbf{B}_n M) = \frac{|\vec{p}_{cm}| \tau_{\mathbf{B}_c}}{8\pi m^2_{\mathbf{B}_c}} |\mathcal{A}(\mathbf{B}_c \to \mathbf{B}_n M)|^2, \qquad (10)
$$

where $|\vec{p}_{cm}| = \sqrt{[(m^2_{\mathbf{B}_c} - (m_{\mathbf{B}_n} + m_M)^2][(m^2_{\mathbf{B}_c} - (m_{\mathbf{B}_n} - m_M)^2]}/(2m_{\mathbf{B}_c})$ and $\tau_{\mathbf{B}_c}$ is the lifetime (the inverse of the total decay width) of $\mathbf{B}_c$. In Eq. (10), the amplitude squared is defined by

$$
\begin{aligned}
&|\mathcal{A}(\mathbf{B}_c \to \mathbf{B}_n M)|^2 \\
&\quad = \frac{(G_F V_{ij} V_{kl})^2}{2} T^\dagger(\mathbf{B}_c \to \mathbf{B}_n M) T(\mathbf{B}_c \to \mathbf{B}_n M). \qquad (11)
\end{aligned}
$$

Note that, since the Lorentz indices have been neglected in the language of the $SU(3)_f$ symmetry, no contractions of the baryon spins are needed, leading to $T^\dagger(\mathbf{B}_c \to \mathbf{B}_n M) = T^*(\mathbf{B}_c \to \mathbf{B}_n M)$.

III. NUMERICAL RESULTS AND DISCUSSIONS

For the numerical analysis, we note that the contributions of the $SU(3)$ parameters (a_4, a_5, a_6, a_7, h') from $H(\overline{15})$ would be neglected based on the following reasons. First, the contributions to the decay branching rates from $H(\overline{15})$ and $H(6)$ lead to a small ratio of $\mathcal{R}(\overline{15}/6) = c^2_+/c^2_- \simeq 17\%$ in terms of $(c_+, c_-) = (0.76, 1.78)$ from the QCD calculation at the scale $\mu = 1$ GeV in the naive dimensional regularization scheme [41,42]. Second, it is pointed out in Ref. [19] that $O^{(\dagger,\prime)}_+$ belong to $H(\overline{15})$ in the group structure and behave as symmetric operators in color indices, whereas the baryon wave functions are totally antisymmetric, such that the mismatch causes the disappearance of

TABLE I. The T amps of the $\mathbf{B}_c \to \mathbf{B}_n M$ decays.

Ξ^0_c	T amp
$\Sigma^+ K^-$	$2(a_2 + \frac{a_4+a_7}{2})$
$\Sigma^0 \bar{K}^0$	$-\sqrt{2}(a_2 + a_3 - \frac{a_6-a_7}{2})$
$\Xi^0 \pi^0$	$-\sqrt{2}(a_1 - a_3 - \frac{a_4-a_5}{2})$
$\Xi^0 \eta$	$\sqrt{2}c\phi(a_1 - a_3 + 2h + \frac{a_4+a_5+2h'}{2})$ $-2s\phi(a_2 + h + \frac{a_7+h'}{2})$
$\Xi^0 \eta'$	$\sqrt{2}s\phi(a_1 - a_3 + 2h + \frac{a_4+a_5+2h'}{2})$ $+c\phi(a_2 + h + \frac{a_7+h'}{2})$
$\Xi^- \pi^+$	$2(a_1 + \frac{a_5+a_6}{2})$
$\Lambda^0 \bar{K}^0$	$-\sqrt{\frac{2}{3}}(2a_1 - a_2 - a_3 + \frac{2a_5-a_6-a_7}{2})$
$\Sigma^+ \pi^-$	$-2(a_2 + \frac{a_4+a_7}{2})s_c$
$\Sigma^- \pi^+$	$-2(a_1 + \frac{a_5+a_6}{2})s_c$
$\Sigma^0 \pi^0$	$-(a_2 + a_3 - \frac{a_4-a_5+a_6-a_7}{2})s_c$
$\Sigma^0 \eta$	$[-c\phi(a_1 + a_2 + 2h + \frac{a_4+a_5-a_6+a_7+2h'}{2})$ $-\sqrt{2}s\phi(a_3 - h - \frac{a_6+h'}{2})]s_c$
$\Sigma^0 \eta'$	$[-s\phi(a_1 + a_2 + 2h + \frac{a_4+a_5-a_6+a_7+2h'}{2})$ $+\sqrt{2}c\phi(a_3 - h - \frac{a_6+h'}{2})]s_c$
$\Xi^- K^+$	$2(a_1 + \frac{a_5+a_6}{2})s_c$
pK^-	$2(a_2 + \frac{a_4+a_7}{2})s_c$
$\Xi^0 K^0$	$2(a_1 - a_2 - a_3 + \frac{a_5-a_7}{2})s_c$
$n\bar{K}^0$	$-2(a_1 - a_2 - a_3 + \frac{a_5-a_7}{2})s_c$
$\Lambda^0 \pi^0$	$\sqrt{\frac{1}{3}}(a_1 + a_2 - 2a_3 + \frac{a_4-a_5-a_6-a_7}{2})s_c$
$\Lambda^0 \eta$	$[\frac{\sqrt{3}c\phi}{3}(a_1 + a_2 - 2a_3 + 6h + \frac{3a_4+a_5+a_6+a_7+6h'}{2})$ $-\frac{\sqrt{6}s\phi}{2}(2a_1 + 2a_2 - a_3 + 3h + \frac{2a_5-a_6+2a_7+3h'}{2})]s_c$
$\Lambda^0 \eta'$	$[\frac{\sqrt{3}s\phi}{3}(a_1 + a_2 - 2a_3 + 6h + \frac{3a_4+a_5+a_6+a_7+6h'}{2})$ $+\frac{\sqrt{6}c\phi}{2}(2a_1 + 2a_2 - a_3 + 3h + \frac{2a_5-a_6+2a_7+3h'}{2})]s_c$
$p\pi^-$	$-2(a_2 + \frac{a_4+a_7}{2})s^2_c$
$\Sigma^- K^+$	$-2(a_1 + \frac{a_5+a_6}{2})s^2_c$
$\Sigma^0 K^0$	$\sqrt{2}(a_1 + \frac{a_5-a_6}{2})s^2_c$
$n\pi^0$	$\sqrt{2}(a_2 - \frac{a_4-a_7}{2})s^2_c$
$n\eta$	$[-\sqrt{2}c\phi(a_2 - 2h + \frac{a_4-a_7-2h'}{2})$ $+2s\phi(a_1 - a_3 + h + \frac{a_5+h'}{2})]s^2_c$
$n\eta'$	$[-\sqrt{2}s\phi(a_2 - 2h + \frac{a_4-a_7-2h'}{2})$ $-2c\phi(a_1 - a_3 + h + \frac{a_5+h'}{2})]s^2_c$
$\Lambda^0 K^0$	$-\sqrt{\frac{2}{3}}(a_1 - 2a_2 - 2a_3 + \frac{a_5+a_6-2a_7}{2})s^2_c$

Ξ^+_c	T amp
$\Sigma^+ \bar{K}^0$	$-2(a_3 - \frac{a_4+a_6}{2})$
$\Xi^0 \pi^+$	$2(a_3 + \frac{a_4+a_6}{2})$
$\Sigma^0 \pi^+$	$\sqrt{2}(a_1 - a_2 + \frac{a_4-a_5+a_6+a_7}{2})s_c$
$\Sigma^+ \pi^0$	$-\sqrt{2}(a_1 - a_2 - \frac{a_4+a_5+a_6-a_7}{2})s_c$
$\Sigma^+ \eta$	$[\sqrt{2}c\phi(a_1 + a_2 + 2h - \frac{a_4+a_5+a_6+a_7-2h'}{2})$ $+2s\phi(a_3 - h - \frac{a_6-h'}{2})]s_c$

(Table continued)

TABLE I. *(Continued)*

Ξ_c^+	T amp
$\Sigma^+\eta'$	$[\sqrt{2}s\phi(a_1+a_2+2h-\frac{a_4+a_5+a_6+a_7-2h'}{2})$ $-2c\phi(a_3-h-\frac{a_6-h'}{2})]s_c$
Ξ^0K^+	$2(a_2+a_3+\frac{a_6-a_7}{2})s_c$
$p\bar{K}^0$	$2(a_1-a_3+\frac{a_4-a_5}{2})s_c$
$\Lambda^0\pi^+$	$\sqrt{\frac{2}{3}}(a_1+a_2-2a_3-\frac{3a_4+a_5+a_6+a_7}{2})s_c$
Σ^0K^+	$\sqrt{2}(a_1-\frac{a_5-a_6}{2})s_c^2$
Σ^+K^0	$2(a_1-\frac{a_5+a_6}{2})s_c^2$
$p\pi^0$	$\sqrt{2}(a_2+\frac{a_4-a_7}{2})s_c^2$
$p\eta$	$[\sqrt{2}c\phi(-a_2+2h+\frac{a_4+a_7+2h'}{2})$ $+2s\phi(a_1-a_3+h-\frac{a_5-2h+h'}{2})]s_c$
$p\eta'$	$[\sqrt{2}s\phi(-a_2+2h+\frac{a_4+a_7+2h'}{2})$ $-2c\phi(a_1-a_3+h-\frac{a_5-2h+h'}{2})]s_c$
$n\pi^+$	$2(a_2-\frac{a_4+a_7}{2})s_c^2$
Λ^0K^+	$\sqrt{\frac{2}{3}}(a_1-2a_2-2a_3-\frac{a_5+a_6-2a_7}{2})s_c^2$

Λ_c^+	T amp
$\Sigma^0\pi^+$	$-\sqrt{2}(a_1-a_2-a_3-\frac{a_5-a_7}{2})$
$\Sigma^+\pi^0$	$\sqrt{2}(a_1-a_2-a_3-\frac{a_5-a_7}{2})$
$\Sigma^+\eta$	$\sqrt{2}c\phi(-a_1-a_2+a_3-2h+\frac{a_5+a_7+2h'}{2})$ $+s\phi(-a_4+2h-h')$
$\Sigma^+\eta'$	$\frac{\sqrt{2}s\phi}{2}(-a_1-a_2+a_3-2h+\frac{a_5+a_7+2h'}{2})$ $-c\phi(-a_4+2h-h')$
Ξ^0K^+	$-2(a_2-\frac{a_4+a_7}{2})$
$p\bar{K}^0$	$-2(a_1-\frac{a_5+a_6}{2})$
$\Lambda^0\pi^+$	$-\sqrt{\frac{2}{3}}(a_1+a_2+a_3-\frac{a_5-2a_6+a_7}{2})$
Σ^+K^0	$-2(a_1-a_3-\frac{a_4-a_5}{2})s_c$
Σ^0K^+	$-\sqrt{2}(a_1-a_3-\frac{a_4+a_5}{2})s_c$
$p\pi^0$	$-\sqrt{2}(a_2+a_3-\frac{a_6-a_7}{2})s_c$
$p\eta$	$[\sqrt{2}c\phi(a_2-a_3+2h+\frac{a_6-a_7-2h'}{2})$ $+2s\phi(-a_1-h+\frac{a_4+a_5+a_6+h'}{2})]s_c$
$p\eta'$	$[\sqrt{2}s\phi(a_2-a_3+2h+\frac{a_6-a_7-2h'}{2})$ $-2c\psi(-a_1-h+\frac{a_4+a_5+a_6+h'}{2})]s_c$
$n\pi^+$	$-2(a_2+a_3-\frac{a_4+a_7}{2})s_c$
Λ^0K^+	$-\sqrt{\frac{2}{3}}(a_1-2a_2+a_3-\frac{3a_4-a_5+2a_6+2a_7}{2})s_c$
pK^0	$2(a_3-\frac{a_4+a_6}{2})s_c^2$
nK^+	$-2(a_3+\frac{a_4+a_6}{2})s_c^2$

$c_+O_+^{(\dagger,\prime)}$ in the calculation of the nonfacotrizable effects, which are regarded to be significant in the charmed baryon decays. Note that, even though the single ignoring of $H(\overline{15})$ is viable, a possible interference between the amplitudes with $H(6)$ and $H(\overline{15})$ may be sizable to fail this assumption, which will be tested in the fit. Hence, being from $H(6)$, the parameters (a_1, a_2, a_3, h) in Eq. (9) are kept for the fit and are, in fact, complex. Since an overall phase can be removed without losing generality, we set a_1 to be real, such that there are seven real independent parameters to be determined, given by

$$a_1, a_2e^{i\delta_{a_2}}, a_3e^{i\delta_{a_3}}, he^{i\delta_h}. \tag{12}$$

We use the minimum χ^2 fit for the determination, given by

$$\chi^2=\sum_i\left(\frac{\mathcal{B}^i_{\rm th}-\mathcal{B}^i_{\rm ex}}{\sigma^i_{\rm ex}}\right)^2+\sum_j\left(\frac{\mathcal{R}^j_{\rm th}-\mathcal{R}^j_{\rm ex}}{\sigma^j_{\rm ex}}\right)^2, \tag{13}$$

where $\mathcal{B}^i_{\rm th}$ and $\mathcal{R}^j_{\rm th}$ stand for the separated decay branching ratios and the ratios of the two-decay branching fractions from the $SU(3)$ amplitudes, while $\mathcal{B}^i_{\rm ex}$ and $\mathcal{R}^j_{\rm ex}$ are the corresponding experimental data, along with $\sigma^i_{\rm ex}$ and $\sigma^j_{\rm ex}$ the 1σ uncertainties, respectively. With the ten experimental data in Table II, the global fit results in

$$\begin{aligned}(a_1,a_2,a_3,h)&=(0.244\pm0.006, 0.115\pm0.014, 0.088\\&\quad\pm0.019, 0.105\pm0.073)\ \mathrm{GeV}^3,\\(\delta_{a_2},\delta_{a_3},\delta_h)&=(78.1\pm7.1, 35.1\pm8.7, 10.2\pm29.6)^\circ,\\\chi^2/\mathrm{d.o.f}&=5.32/3=1.77,\end{aligned} \tag{14}$$

where d.o.f represents the degree of freedom. The numerical values for the parameters in Eq. (14) are the theoretical inputs, which are used to predict the two-body $\mathbf{B}_c\to\mathbf{B}$ decays in Table III.

Since the value of $\chi^2/\mathrm{d.o.f}\simeq1.8$ in Eq. (14) indicates a good fit, there exists no inconstancy by neglecting $H(\overline{15})$ in our analysis. Note that the determinations of $|a_1|$ and $|a_2|$ depend on $T(\Lambda_c^+\to p\bar{K}^0)=-2a_1$ and $T(\Lambda_c^+\to\Xi^0K^+)=-2a_2$ in Table I, respectively, by ignoring (a_5+a_6) and (a_4+a_7), associated with $H(\overline{15})$. Similarly, one can extract $|a_3|$ based on $T(\Xi_c^+\to\Xi^0\pi^+)=2a_3+(a_4+a_6)\simeq2a_3$. Consequently, we get

TABLE II. The data of the $\mathbf{B}_c\to\mathbf{B}_nM$ decays.

Branching ratios	Data [4,7]
$10^2\mathcal{B}(\Lambda_c^+\to p\bar{K}^0)$	3.16 ± 0.16
$10^2\mathcal{B}(\Lambda_c^+\to\Lambda\pi^+)$	1.30 ± 0.07
$10^2\mathcal{B}(\Lambda_c^+\to\Sigma^+\pi^0)$	1.24 ± 0.10
$10^2\mathcal{B}(\Lambda_c^+\to\Sigma^0\pi^+)$	1.29 ± 0.07
$10^2\mathcal{B}(\Lambda_c^+\to\Xi^0K^+)$	0.50 ± 0.12
$10^2\mathcal{B}(\Lambda_c^+\to\Sigma^+\eta)$	0.70 ± 0.23
$10^4\mathcal{B}(\Lambda_c^+\to\Lambda K^+)$	6.1 ± 1.2
$10^4\mathcal{B}(\Lambda_c^+\to\Sigma^0K^+)$	5.2 ± 0.8
$10^4\mathcal{B}(\Lambda_c^+\to p\eta)$	12.4 ± 3.0
$\mathcal{R}=\frac{\mathcal{B}(\Xi_c^0\to\Lambda\bar{K}^0)}{\mathcal{B}(\Xi_c^0\to\Xi^-\pi^+)}$	0.420 ± 0.056

TABLE III. The numerical results of the $\mathbf{B}_c \to \mathbf{B}_n M$ decays with $\mathcal{B}_{\mathbf{B}_n M} \equiv \mathcal{B}(\mathbf{B}_c \to \mathbf{B}_n M)$, where the number with the dagger (†) is the reproduction of the experimental data input, instead of the prediction.

Ξ_c^0	Our results	Reference [43]
$10^3\mathcal{B}_{\Sigma^+K^-}$	3.5 ± 0.9	3.1 ± 0.9
$10^3\mathcal{B}_{\Sigma^0\bar{K}^0}$	4.7 ± 1.2	4.6 ± 1.4
$10^3\mathcal{B}_{\Xi^0\pi^0}$	4.3 ± 0.9	0.7–18.1
$10^3\mathcal{B}_{\Xi^0\eta}$	$1.7^{+1.0}_{-1.7}$	
$10^3\mathcal{B}_{\Xi^0\eta'}$	$8.6^{+11.0}_{-6.3}$	
$10^3\mathcal{B}_{\Xi^-\pi^+}$	15.7 ± 0.7	22.4 ± 3.4
$10^3\mathcal{B}_{\Lambda^0\bar{K}^0}$	8.3 ± 0.9	9.4 ± 1.6
$10^4\mathcal{B}_{\Sigma^+\pi^-}$	2.0 ± 0.5	
$10^4\mathcal{B}_{\Sigma^-\pi^+}$	9.0 ± 0.4	
$10^4\mathcal{B}_{\Sigma^0\pi^0}$	3.2 ± 0.3	
$10^4\mathcal{B}_{\Sigma^0\eta}$	$3.6^{+1.0}_{-0.9}$	
$10^4\mathcal{B}_{\Sigma^0\eta'}$	$1.7^{+3.0}_{-1.5}$	
$10^4\mathcal{B}_{\Xi^-K^+}$	7.6 ± 0.4	
$10^4\mathcal{B}_{\Xi^0K^0}$	6.3 ± 1.2	
$10^4\mathcal{B}_{pK^-}$	2.1 ± 0.5	
$10^4\mathcal{B}_{n\bar{K}^0}$	7.9 ± 1.4	
$10^4\mathcal{B}_{\Lambda^0\pi^0}$	0.2 ± 0.2	
$10^4\mathcal{B}_{\Lambda^0\eta}$	$1.6^{+1.2}_{-0.8}$	
$10^4\mathcal{B}_{\Lambda^0\eta'}$	$9.4^{+11.6}_{-6.8}$	
$10^6\mathcal{B}_{p\pi^-}$	12.1 ± 3.1	
$10^6\mathcal{B}_{\Sigma^-K^+}$	44.5 ± 2.1	
$10^6\mathcal{B}_{\Sigma^0K^0}$	22.3 ± 1.0	
$10^6\mathcal{B}_{n\pi^0}$	6.0 ± 1.5	
$10^6\mathcal{B}_{n\eta}$	$26.5^{+11.4}_{-10.1}$	
$10^6\mathcal{B}_{n\eta'}$	$30.7^{+42.3}_{-24.4}$	
$10^6\mathcal{B}_{\Lambda^0K^0}$	14.4 ± 3.7	

Ξ_c^+	Our results	Reference [43]
$10^3\mathcal{B}_{\Sigma^+\bar{K}^0}$	8.0 ± 3.9	0.1–102.2
$10^3\mathcal{B}_{\Xi^0\pi^+}$	8.1 ± 4.0	1.2–96.8
$10^4\mathcal{B}_{\Sigma^0\pi^+}$	18.5 ± 2.2	
$10^4\mathcal{B}_{\Sigma^+\pi^0}$	18.5 ± 2.2	
$10^4\mathcal{B}_{\Sigma^+\eta}$	$28.4^{+8.2}_{-6.9}$	
$10^4\mathcal{B}_{\Sigma^+\eta'}$	$13.2^{+24.0}_{-11.9}$	
$10^4\mathcal{B}_{\Xi^0K^+}$	18.0 ± 4.7	
$10^4\mathcal{B}_{p\bar{K}^0}$	20.3 ± 4.2	
$10^4\mathcal{B}_{\Lambda^0\pi^+}$	1.6 ± 1.2	
$10^5\mathcal{B}_{\Sigma^0K^+}$	8.8 ± 0.4	
$10^5\mathcal{B}_{\Sigma^+K^0}$	17.6 ± 0.8	
$10^6\mathcal{B}_{p\pi^0}$	23.8 ± 6.1	
$10^5\mathcal{B}_{p\eta}$	$10.5^{+4.5}_{-4.0}$	

(Table continued)

TABLE III. *(Continued)*

Ξ_c^+	Our results	Reference [43]
$10^5\mathcal{B}_{p\eta'}$	$12.1^{+16.7}_{-9.7}$	
$10^6\mathcal{B}_{n\pi^+}$	47.6 ± 12.2	
$10^6\mathcal{B}_{\Lambda^0K^+}$	56.8 ± 14.5	

Λ_c^+	Our results	Reference [43]
$10^3\mathcal{B}_{\Sigma^0\pi^+}$	$(1.3 \pm 0.2)^\dagger$	$(1.27 \pm 0.17)^\dagger$
$10^3\mathcal{B}_{\Sigma^+\pi^0}$	$(1.3 \pm 0.2)^\dagger$	$(1.27 \pm 0.17)^\dagger$
$10^2\mathcal{B}_{\Sigma^+\eta}$	$(0.7^{+0.4}_{-0.3})^\dagger$	
$10^2\mathcal{B}_{\Sigma^+\eta'}$	$1.0^{+1.6}_{-0.8}$	
$10^2\mathcal{B}_{\Xi^0K^+}$	$(0.5 \pm 0.1)^\dagger$	$(0.50 \pm 0.12)^\dagger$
$10^2\mathcal{B}_{p\bar{K}^0}$	$(3.3 \pm 0.2)^\dagger$	$(2.72 - 3.60)^\dagger$
$10^2\mathcal{B}_{\Lambda^0\pi^+}$	$(1.3 \pm 0.2)^\dagger$	$(1.30 \pm 0.17)^\dagger$
$10^4\mathcal{B}_{\Sigma^+K^0}$	8.0 ± 1.6	
$10^4\mathcal{B}_{\Sigma^0K^+}$	$(4.0 \pm 0.8)^\dagger$	
$10^4\mathcal{B}_{p\pi^0}$	5.7 ± 1.5	
$10^4\mathcal{B}_{p\eta}$	$(12.5^{+3.8}_{-3.6})^\dagger$	
$10^4\mathcal{B}_{p\eta'}$	$12.2^{+14.3}_{-8.7}$	
$10^4\mathcal{B}_{n\pi^+}$	11.3 ± 2.9	
$10^4\mathcal{B}_{\Lambda^0K^+}$	$(4.6 \pm 0.9)^\dagger$	
$10^6\mathcal{B}_{pK^0}$	12.2 ± 6.0	
$10^6\mathcal{B}_{nK^+}$	12.2 ± 6.0	

$$R_0\mathcal{B}(\Lambda_c^+ \to p\bar{K}^0) = \mathcal{B}(\Xi_c^0 \to \Xi^-\pi^+) = (15.7 \pm 0.7) \times 10^{-3},$$
$$R_0\mathcal{B}(\Lambda_c^+ \to \Xi^0K^+) = \mathcal{B}(\Xi_c^0 \to \Sigma^+K^-) = (0.4 \pm 0.1) \times 10^{-2},$$
$$\mathcal{B}(\Xi_c^+ \to \Sigma^+\bar{K}^0) = \mathcal{B}(\Xi_c^+ \to \Xi^0\pi^+) = (8.1 \pm 4.0) \times 10^{-3}, \tag{15}$$

without the contributions from $H(\overline{15})$, where $R_0 = \tau_{\Xi_c^0}/\tau_{\Lambda_c^+} = 0.56 \pm 0.07$. To check if the $H(\overline{15})$ terms are indeed negligible, we may use the relations from Table I, given by

$$T(\Lambda_c^+ \to p\bar{K}^0) + T(\Xi_c^0 \to \Xi^-\pi^+) = 2(a_5 + a_6),$$
$$T(\Lambda_c^+ \to \Xi^0K^+) + T(\Xi_c^0 \to \Sigma^+K^-) = 2(a_4 + a_7),$$
$$T(\Xi_c^+ \to \Xi^0\pi^+) + T(\Xi_c^+ \to \Sigma^+\bar{K}^0) = 2(a_4 + a_6). \tag{16}$$

Clearly, if the results in Eq. (15) do not agree with the future measurements, the contributions from $H(\overline{15})$ should be reconsidered as seen in Eq. (16).

According to the PDG [7], the branching fractions in the Ξ_c^0 decays are observed to be relative to $\mathcal{B}_{\Xi^-\pi^+} \equiv \mathcal{B}(\Xi_c^0 \to \Xi^-\pi^+)$, predicted in Table III. Hence, by using the partial observations of $\mathcal{B}(\Xi_c^0 \to \Lambda K^-\pi^+) = (1.07 \pm 0.14)\mathcal{B}_{\Xi^-\pi^+}$, $\mathcal{B}(\Xi_c^0 \to \Lambda K^+K^-) = (0.029 \pm 0.007)\mathcal{B}_{\Xi^-\pi^+}$, and $\mathcal{B}(\Xi_c^0 \to \Xi^-e^+\nu_e) = (3.1 \pm 1.1)\mathcal{B}_{\Xi^-\pi^+}$, we obtain

$$\mathcal{B}(\Xi_c^0 \to \Lambda K^- \pi^+) = (16.8 \pm 2.3) \times 10^{-3},$$
$$\mathcal{B}(\Xi_c^0 \to \Lambda K^+ K^-) = (4.5 \pm 1.1) \times 10^{-4},$$
$$\mathcal{B}(\Xi_c^0 \to \Xi^- e^+ \nu_e) = (48.7 \pm 17.4) \times 10^{-3}. \quad (17)$$

Similarly, the branching fractions in the Ξ_c^+ decays are measured to be relative to $\mathcal{B}(\Xi_c^+ \to \Xi^- \pi^+ \pi^+)$, which has not been theoretically and experimentally studied yet. With $\mathcal{B}(\Xi_c^+ \to \Xi^0 \pi^+)/\mathcal{B}(\Xi_c^+ \to \Xi^- \pi^+ \pi^+) = 0.55 \pm 0.16$ [7] and $\mathcal{B}(\Xi_c^+ \to \Xi^0 \pi^+)$ in Table III, we find

$$\mathcal{B}_{\Xi^- 2\pi^+} \equiv \mathcal{B}(\Xi_c^+ \to \Xi^- \pi^+ \pi^+) = (14.7 \pm 8.4) \times 10^{-3}. \quad (18)$$

Subsequently, the relative branching fractions of $\mathcal{B}(\Xi_c^+ \to pK_s^0 K_s^0) = (0.087 \pm 0.021)\mathcal{B}_{\Xi^- 2\pi^+}$, $\mathcal{B}(\Xi_c^+ \to \Sigma^+ K^- \pi^+) = (0.94 \pm 0.10)\mathcal{B}_{\Xi^- 2\pi^+}$, $\mathcal{B}(\Xi_c^+ \to \Xi^0 \pi^+ \pi^0) = (2.3 \pm 0.7)\mathcal{B}_{\Xi^- 2\pi^+}$ and $\mathcal{B}(\Xi_c^+ \to \Xi^0 e^+ \nu_e) = (2.3^{+0.7}_{-0.8})\mathcal{B}_{\Xi^- 2\pi^+}$ [7] lead to

$$\mathcal{B}(\Xi_c^+ \to pK_s^0 K_s^0) = (1.3 \pm 0.8) \times 10^{-3},$$
$$\mathcal{B}(\Xi_c^+ \to \Sigma^+ K^- \pi^+) = (13.8 \pm 8.0) \times 10^{-3},$$
$$\mathcal{B}(\Xi_c^+ \to \Xi^0 \pi^+ \pi^0) = (33.8 \pm 21.9) \times 10^{-3},$$
$$\mathcal{B}(\Xi_c^+ \to \Xi^0 e^+ \nu_e) = (33.8^{+21.9}_{-22.6}) \times 10^{-3}. \quad (19)$$

By adding the $h^{(\prime)}$ terms, we are able to include the contributions from the singlet η_1 in the $SU(3)_f$ amplitudes, which have been used to explain the observations of $\mathcal{B}(\Lambda_c^+ \to \Sigma^+ \eta)$ and $\mathcal{B}(\Lambda_c^+ \to p\eta)$. Nonetheless, the estimations of $\mathcal{B}(\Lambda_c^+ \to \Sigma^+(p)\eta') \simeq \mathcal{B}(\Lambda_c^+ \to \Sigma^+(p)\eta)$ [30] show no inequality as $\mathcal{B}(B \to K\eta') \gg \mathcal{B}(B \to K\eta)$ or $\mathcal{B}(B \to K^*\eta) \gg \mathcal{B}(B \to K^*\eta')$. On the other hand, it is interesting to note that, despite of the large uncertainties, the $\Xi_c \to \mathbf{B}_n \eta^{(\prime)}$ decays contain the similar inequalities between the η and η' modes, given by

$$\mathcal{B}(\Xi_c^0 \to \Xi^0 \eta, \Xi^0 \eta') = (1.7^{+1.0}_{-1.7}, 8.6^{+11.0}_{-6.3}) \times 10^{-3},$$
$$\mathcal{B}(\Xi_c^0 \to \Lambda^0 \eta, \Lambda^0 \eta') = (1.6^{+1.2}_{-0.8}, 9.4^{+11.6}_{-6.8}) \times 10^{-4},$$
$$\mathcal{B}(\Xi_c^+ \to \Sigma^+ \eta, \Sigma^+ \eta') = (28.4^{+8.2}_{-6.9}, 13.2^{+24.0}_{-11.9}) \times 10^{-4}. \quad (20)$$

We remark that, as shown in Table III, our numerical results for the Cabibbo-allowed processes are consistent with those in Ref. [43], where $\mathcal{B}(\mathbf{B}_c \to \mathbf{B}_n \bar{K}^0)$ are taken from $\mathcal{B}(\mathbf{B}_c \to \mathbf{B}_n K_S^0)$. Finally, we emphasize that there is a discrepancy between the theory and data for $\mathcal{B}(\Lambda_c^+ \to p\pi^0)$. In Table III, $\mathcal{B}(\Lambda_c^+ \to p\pi^0)$ is predicted to be $(5.7 \pm 1.5) \times 10^{-4}$, whereas it is measured to be less than 3×10^{-4} [4]. Nonetheless, the estimation in the factorization approach also gives $\mathcal{B}(\Lambda_c^+ \to p\pi^0) = f_\pi^2/(2f_K^2)s_c^2\mathcal{B}(\Lambda_c^+ \to p\bar{K}^0) = (5.5 \pm 0.3) \times 10^{-4}$ to be as large as our $SU(3)_f$ prediction in Table III, with the experimental input of $\mathcal{B}(\Lambda_c^+ \to p\bar{K}^0) = (3.16 \pm 0.16) \times 10^{-2}$ [7]. Clearly, to resolve this inconsistency, it is necessary to remeasure the decay of $\Lambda_c^+ \to p\pi^0$ in the future experiment.

IV. CONCLUSION

With the $SU(3)_f$ symmetry, we have studied the two-body antitriplet charmed baryon weak decays. We have predicted that $\mathcal{B}(\Xi_c^0 \to \Xi^- \pi^+) = (15.7 \pm 0.7) \times 10^{-3}$ and $\mathcal{B}(\Xi_c^+ \to \Xi^- \pi^+ \pi^+) = (14.7 \pm 8.4) \times 10^{-3}$, while the branching ratios of the Ξ_c^0 and Ξ_c^+ decays are measured to be relative to $\mathcal{B}(\Xi_c^0 \to \Xi^- \pi^+)$ and $\mathcal{B}(\Xi_c^+ \to \Xi^- \pi^+ \pi^+)$, respectively. Hence, we have extracted that $\mathcal{B}(\Xi_c^0 \to \Lambda K^- \pi^+, \Lambda K^+ K^-, \Xi^- e^+ \nu_e) = (16.8 \pm 2.3, 0.45 \pm 0.11, 48.7 \pm 17.4) \times 10^{-3}$ and $\mathcal{B}(\Xi_c^+ \to pK_s^0 K_s^0, \Sigma^+ K^- \pi^+, \Xi^0 \pi^+ \pi^0, \Xi^0 e^+ \nu_e) = (1.3 \pm 0.8, 13.8 \pm 8.0, 33.8 \pm 21.9, 33.8^{+21.9}_{-22.6}) \times 10^{-3}$. In addition, we have shown that $\mathcal{B}(\Xi_c^0 \to \Xi^0 \eta, \Xi^0 \eta') = (1.7^{+1.0}_{-1.7}, 8.6^{+11.0}_{-6.3}) \times 10^{-3}$, $\mathcal{B}(\Xi_c^0 \to \Lambda^0 \eta, \Lambda^0 \eta') = (1.6^{+1.2}_{-0.8}, 9.4^{+11.6}_{-6.8}) \times 10^{-4}$, and $\mathcal{B}(\Xi_c^+ \to \Sigma^+ \eta, \Sigma^+ \eta') = (28.4^{+8.2}_{-6.9}, 13.2^{+24.0}_{-11.9}) \times 10^{-4}$, representing the inequalities, similar to those of $\mathcal{B}(B \to K\eta') \gg \mathcal{B}(B \to K\eta)$ or $\mathcal{B}(B \to K^*\eta) \gg \mathcal{B}(B \to K^*\eta')$ in the mesonic B decays involving $\eta^{(\prime)}$. According to our predictions, the branching ratios of two and three-body Ξ_c decays are accessible to the experiments at BESIII and LHCb.

ACKNOWLEDGMENTS

This work was supported in part by National Center for Theoretical Sciences, MoST (Grant No. MoST-104-2112-M-007-003-MY3) and National Science Foundation of China (Grant No. 11675030).

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