Doubly charmful baryonic *B* decays

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There are two apparent puzzles connected with the two-body and three-body doubly charmed baryonic *B* decays. First, earlier calculations based on QCD sum rules or the diquark model predict $\mathcal{B}(\bar{B}^0 \to \Xi_c^+ \bar{\Lambda}_c^-) \approx \mathcal{B}(\bar{B}^0 \to \mathcal{B}_c \bar{N})$, while experimentally the former has a rate 2 orders of magnitude larger than the latter. Second, a naive estimate of the branching ratio $\mathcal{O}(10^{-9})$ for the color-suppressed three-body decay $\bar{B} \to \Lambda_c^+ \bar{\Lambda}_c^- K$, which is highly suppressed by phase space, is too small by 5 to 6 orders of magnitude compared to the experiment. We show that the great suppression for the $\Lambda_c^+ \bar{\Lambda}_c^- K$ production can be alleviated provided that there exists a narrow hidden charm bound state with a mass near the $\Lambda_c \bar{\Lambda}_c$ threshold. This new state that couples strongly to the charmed baryon pair can be searched for in *B* decays and in $p\bar{p}$ collisions by studying the mass spectrum of $D^{(*)}\bar{D}^{(*)}$ or $\Lambda_c \bar{\Lambda}_c$. The doubly charmful decay $\bar{B} \to \Xi_c \bar{\Lambda}_c$ has a configuration more favorable than the singly charmful one such as $\bar{B}^0 \to \Lambda_c \bar{p}$ since no hard gluon is needed to produce the energetic $\Xi_c \bar{\Lambda}_c$ pair in the former decay, while two hard gluons are needed for the latter process. Assuming that a soft $q\bar{q}$ quark pair is produced through the σ and π meson exchanges in the configuration for $\bar{B} \to \Xi_c \bar{\Lambda}_c$, it is found that its branching ratio is of order 10^{-3}, in agreement with the experiment.

I. INTRODUCTION

Recently Belle has observed for the first time two-body and three-body doubly charmed baryonic *B* decays in which two charmed baryons are produced in the final state [1,2]. The measured branching ratios are

$$\mathcal{B}(B^- \to \Lambda_c^+ \bar{\Lambda}_c^- K^-) = (6.5^{+1.0}_{-0.9} \pm 0.8 \pm 3.4) \times 10^{-4},$$
$$\mathcal{B}(B^0 \to \Lambda_c^+ \bar{\Lambda}_c^- K^0) = (7.9^{+2.9}_{-2.3} \pm 1.2 \pm 4.2) \times 10^{-4},$$

(1)

for three-body decays and

$$\mathcal{B}(B^- \to \Xi_c^0 \bar{\Lambda}_c^-)\mathcal{B}(\Xi_c^0 \to \Xi^- \pi^+)$$
$$= (4.8^{+1.0}_{-0.9} \pm 1.1 \pm 1.2) \times 10^{-5},$$
$$\mathcal{B}(\bar{B}^0 \to \Xi_c^+ \bar{\Lambda}_c^-)\mathcal{B}(\Xi_c^+ \to \Xi^- \pi^+ \pi^+)$$
$$= (9.3^{+3.7}_{-2.8} \pm 1.9 \pm 2.4) \times 10^{-5}$$

(2)

for two-body decays. Taking the theoretical estimates (see e.g. Table III of [3]), $\mathcal{B}(\Xi_c^0 \to \Xi^- \pi^+) \approx 1.3\%$ and $\mathcal{B}(\Xi_c^+ \to \Xi^0 \pi^+) \approx 3.9\%$ together with the experimental measurement $\mathcal{B}(\Xi_c^+ \to \Xi^0 \pi^+)/\mathcal{B}(\Xi_c^+ \to \Xi^- \pi^+ \pi^+) = 0.55 \pm 0.16$ [4], it follows that

$$\mathcal{B}(B^- \to \Xi_c^0 \bar{\Lambda}_c^-) \approx 4.8 \times 10^{-3},$$
$$\mathcal{B}(\bar{B}^0 \to \Xi_c^+ \bar{\Lambda}_c^-) \approx 1.2 \times 10^{-3}.$$

(3)

Therefore, the two-body doubly charmed baryonic *B* decay $B \to \mathcal{B}_c \bar{\mathcal{B}}_c'$ has a branching ratio of order 10^{-3}, to be compared with [5,6]

$$\mathcal{B}(\bar{B}^0 \to \Lambda_c^+ \bar{p}) = (2.19^{+0.56}_{-0.49} \pm 0.32 \pm 0.57) \times 10^{-5},$$
$$\mathcal{B}(B^- \to \Sigma_c(2455)^0 \bar{p}) = (3.67^{+0.74}_{-0.66} \pm 0.36 \pm 0.95) \times 10^{-5},$$

(4)

for singly charmed baryonic *B* decays and [7,8]

$$\mathcal{B}(B^0 \to p\bar{p}) < 2.7 \times 10^{-7},$$
$$\mathcal{B}(B^0 \to \Lambda\bar{\Lambda}) < 6.9 \times 10^{-7},$$
$$\mathcal{B}(B^- \to \Lambda\bar{p}) < 4.9 \times 10^{-7},$$

(5)

for charmless baryonic *B* decays. Therefore, we have the pattern

$$\mathcal{B}_c \bar{\mathcal{B}}_c'(\sim 10^{-3}) \gg \mathcal{B}_c \bar{\mathcal{B}}(\sim 10^{-5}) \gg \mathcal{B}_1 \bar{\mathcal{B}}_2(\sim 10^{-7}),$$

(6)

for two-body baryonic *B* decays.

Using $\mathcal{B}(\bar{B}^0 \to \Lambda_c^+ \bar{p})$ as a benchmark, one will expect a branching ratio of order 10^{-7} for the charmless decay $B \to \mathcal{B}_1 \bar{\mathcal{B}}_2$ after replacing the quark mixing angle V_{cb} by V_{ub}, provided that the dynamical suppression for the latter is neglected. However, since the doubly charmed baryonic decay mode $\Xi_c \bar{\Lambda}_c$ proceeds via $b \to cs\bar{c}$, while $\Lambda_c \bar{p}$ via a $b \to cd\bar{u}$ quark transition, the Cabibbo-Kobayashi-Maskawa (CKM) mixing angles for them are the same in magnitude but opposite in sign. One may wonder why the $\mathcal{B}_c \bar{\mathcal{B}}_c'$ mode has a rate 2 orders of magnitude larger than $\mathcal{B}_c \bar{\mathcal{B}}$. Indeed, earlier calculations based on QCD sum rules [9] or the diquark model [10] all predict that $\mathcal{B}(B \to \Xi_c \bar{\Lambda}_c) \approx \mathcal{B}(\bar{B} \to \mathcal{B}_c \bar{N})$, which is in violent disagreement with the experiment. This implies that some important dynamical suppression effect for the $\mathcal{B}_c \bar{N}$ production with respect to $\Xi_c \bar{\Lambda}_c$ is missing in previous studies.

As for the three-body decay $B \to \Lambda_c \bar{\Lambda}_c K$, its branching ratio is estimated to be of order 10^{-9}, which is extremely small due to the tiny phase space available for this decay and the color-suppression effect. The puzzle is that why the measured rate is much larger than the naive expectation?

A crucial ingredient for understanding the baryonic B decays is the threshold or low-mass enhancement behavior of the baryon-pair invariant mass in the spectrum for $B \to \mathcal{B}_1\bar{\mathcal{B}}_2 M$: It sharply peaks at very low values. That is, the B meson is preferred to decay into a baryon-antibaryon pair with low invariant mass accompanied by a fast recoiled meson. Therefore, some three-body final states have rates larger than their two-body counterparts, e.g. $p\bar{p}K^\pm \gg p\bar{p}$, $\Lambda\bar{p}\pi^\pm \gg \Lambda\bar{p}$, $\Sigma_c\bar{p}\pi^\pm \gg \Sigma_c\bar{p}$.[1] This phenomenon can be understood in terms of the threshold effect, namely, the invariant mass of the dibaryon is preferred to be close to the threshold. The configuration of the two-body decay $B \to \mathcal{B}_1\bar{\mathcal{B}}_2$ is not favorable since its invariant mass is m_B. In $B \to \mathcal{B}_1\bar{\mathcal{B}}_2 M$ decays, the effective mass of the baryon pair is reduced as the emitted meson can carry away a large amount of energies. The two-body decay pattern (6) also follows from the low-mass enhancement effect: The energy release is least for the B decay into two charmed baryons and becomes very large when the final-state baryons are charmless.

Although the gross feature of the baryonic B decays can be qualitatively comprehended in terms of the near threshold effect, how to quantitatively evaluate their absolute decay rates and how to realize the low-mass enhancement effect require detailed dynamical studies. In the present work we will focus on the doubly charmful baryonic B decays, namely, $\bar{B} \to \Xi_c\bar{\Lambda}_c$ and $\bar{B} \to \Lambda_c\bar{\Lambda}_cK$ in Secs. II and III, respectively, aiming to resolve the aforementioned two puzzles connected with them. Section IV gives the conclusion. The evaluation of the delta functions occurring in the phase space integral is discussed in the appendix.

II. THREE-BODY DECAYS

We consider the decay $B^- \to \Lambda_c^+\bar{\Lambda}_c^-K^-$, which proceeds through the internal W-emission diagram in Fig. 1(a). It turns out this diagram is factorizable. In the weak Hamiltonian approach, the factorizable amplitude reads

$$A(B^- \to \Lambda_c^+\bar{\Lambda}_c^-K^-)$$
$$= \frac{G_F}{\sqrt{2}} V_{cs}V_{cb}^* a_2 \langle \Lambda_c^+\bar{\Lambda}_c^-|(\bar{c}c)|0\rangle\langle K^-|(\bar{s}b)|B^-\rangle, \quad (7)$$

where $(\bar{q}_1q_2) \equiv \bar{q}_1\gamma_\mu(1 - \gamma_5)q_2$, and the effective Wilson coefficient a_2 indicates that this decay is color suppressed. The matrix elements can be parametrized as

$$\langle \Lambda_c(p_1)\bar{\Lambda}_c(p_2)|(\bar{c}c)|0\rangle = \bar{u}_{\Lambda_c}(p_1)\left[f_1(q^2)\gamma_\mu + i\frac{f_2(q^2)}{2m_{\Lambda_c}}\sigma_{\mu\nu}q^\nu - \left(g_1(q^2)\gamma_\mu + \frac{g_3(q^2)}{2m_{\Lambda_c}}q_\mu\right)\gamma_5\right]v_{\bar{\Lambda}_c}(p_2),$$

$$\langle K^-(p_K)|(\bar{s}b)|B^-(p_B)\rangle = F_1^{BK}(q^2)(p_B + p_K)_\mu + (F_0^{BK}(q^2) - F_1^{BK}(q^2))\frac{m_B^2 - m_K^2}{q^2}q_\mu, \quad (8)$$

with $q = p_B - p_K = p_1 + p_2$. In terms of the form factors, the decay amplitude has the expression

$$A(B^- \to \Lambda_c^+\bar{\Lambda}_c^-K^-) = \frac{G_F}{\sqrt{2}}V_{cs}V_{cb}^* a_2\bar{u}_{\Lambda_c}[a\slashed{p}_K + b - (c\slashed{p}_K + d)\gamma_5]v_{\bar{\Lambda}_c}, \quad (9)$$

with

$$a = 2F_1^{BK}(q^2)[f_1(q^2) + f_2(q^2)],$$

$$b = 2F_1^{BK}(q^2)f_2(q^2)(p_2 - p_1)\cdot p_K/(2m_{\Lambda_c}),$$

$$c = 2F_1^{BK}(q^2)g_1(q^2),$$

$$d = 2m_{\Lambda_c}g_1(q^2)\left[F_1^{BK}(q^2) + (F_0^{BK}(q^2) - F_1^{BK}(q^2))\frac{m_B^2 - m_K^2}{q^2}\right]$$
$$+ g_3(q^2)F_0^{BK}(q^2)(m_B^2 - m_K^2)/(2m_{\Lambda_c}). \quad (10)$$

There are numerous estimates of the $B \to K$ transition form factors. We will follow [11] where the form factors are evaluated using the relativistic covariant light-front quark model.

Because of the heavy mass of Λ_c, the phase space of the $\Lambda_c^+\bar{\Lambda}_c^-K^-$ decay is about a hundred times smaller than, say, that of the $\Lambda\bar{p}\pi^+$ [12]. The $\Lambda_c^+\bar{\Lambda}_c^-$ form factors, if any, can therefore be taken as constants whose values are determined at the threshold over the phase space. To achieve a rate that is at least comparable with that of the $\Lambda\bar{p}\pi^+$ whose branching ratio is of order 3×10^{-6} [12], one would need the $\Lambda_c^+\bar{\Lambda}_c^-$ form factors to be more than 100 times larger than those of the $\Lambda\bar{p}\pi^+$ near the $\Lambda\bar{p}$ threshold, which is quite unlikely since the $\Lambda\bar{p}$ form
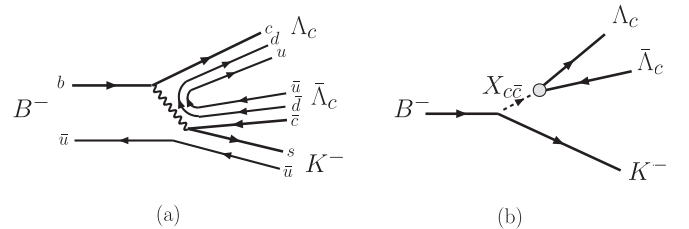
FIG. 1. $B^- \to \Lambda_c^+\bar{\Lambda}_c^-K^-$ as proceeding through (a) the internal W-emission diagram, and (b) the dominant charmoniumlike resonance $X_{c\bar{c}}$. The blob in (b) shows where the strong decays take place.

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[1]The three-body decay is usually referred to as the nonresonant one. The relation $\Lambda_c\bar{p}\pi^- \gg \Lambda_c\bar{p}$ is trivial as the former arises mostly from resonant contributions [6].

factors have their maximum values already about $\mathcal{O}(1)$. Besides, $\mathcal{O}(10^2)$ form factors would just give a rate of $\Lambda_c^+ \bar{\Lambda}_c^- K^-$ comparable to $\sim \mathcal{O}(10^{-6})$, not to mention the remaining factor of $\mathcal{O}(10^2)$ difference between the rates of $\Lambda \bar{p} \pi^+$ and of $\Lambda_c^+ \bar{\Lambda}_c^- K^-$. Therefore we conclude that the suppression from the $\Lambda_c^+ \bar{\Lambda}_c^- K^-$ phase space is so strong that $\Lambda_c^+ \bar{\Lambda}_c^-$ pair is unlikely to be produced dominantly through the direct three-body decay processes. The great suppression, however, seems to hint strongly that a $c\bar{c}$-content resonance with the width comparable to the

nearby resonances like $\psi(4415)$ could be located around the threshold of $\Lambda_c^+ \bar{\Lambda}_c^-$ (~ 4.6 GeV), and the whole process takes place dominantly via the charmoniumlike resonance as shown in Fig. 1(b).

Let us assume the resonance $X_{c\bar{c}}$ exists with a mass $m_{X_{c\bar{c}}} \gtrsim 2m_{\Lambda_c}$ and a width $\Gamma_{X_{c\bar{c}}}$. Let us further assume that this resonance is a spin-1 particle with $J^P = 1^-$ or 1^+, as inspired from the observation that all the charmonia near the $\Lambda_c^+ \bar{\Lambda}_c^-$ threshold are spin-1 particles. The decay amplitude then reads

$$A_{X_{c\bar{c}}}(B^- \to \Lambda_c^+ \bar{\Lambda}_c^- K^-) = \frac{G_F}{\sqrt{2}} V_{cs} V_{cb}^* a_2 \langle K^- | \bar{s} \gamma_\mu (1 - \gamma_5) b | B^- \rangle m_{X_{c\bar{c}}} f_{X_{c\bar{c}}} \left(\frac{-g^{\mu\nu} + \frac{q^\mu q^\nu}{m_{X_{c\bar{c}}}^2}}{q^2 - m_{X_{c\bar{c}}}^2 + i m_{X_{c\bar{c}}} \Gamma_{X_{c\bar{c}}}} \right) i \bar{u}_{\Lambda_c}(p_1) M_\nu v_{\bar{\Lambda}_c}(p_2), \quad (11)$$

where

$$M_\nu = h_1^{\Lambda_c \bar{\Lambda}_c V} \gamma_\nu + \frac{i h_2^{\Lambda_c \bar{\Lambda}_c V}}{2m_{\Lambda_c}} \sigma_{\nu\rho} q^\rho, \quad (12)$$

when $X_{c\bar{c}}$ is a vector ($X_{c\bar{c}} = V$), and

$$M_\nu = \left(h_1^{\Lambda_c \bar{\Lambda}_c A} \gamma_\nu + \frac{h_2^{\Lambda_c \bar{\Lambda}_c A}}{2m_{\Lambda_c}} q_\nu \right) \gamma_5, \quad (13)$$

when $X_{c\bar{c}}$ is an axial-vector ($X_{c\bar{c}} = A$) particle. $f_{X_{c\bar{c}}}$ is the decay constant for $X_{c\bar{c}}$, and $h_{1,2}^{\Lambda_c \bar{\Lambda}_c V}$ and $h_{1,2}^{\Lambda_c \bar{\Lambda}_c A}$ represent the *dimensionless* $\Lambda_c \bar{\Lambda}_c X_{c\bar{c}}$ strong couplings. Since the allowed phase space is very small, the strong couplings can effectively be treated as constants within this region. The decay constant $f_{X_{c\bar{c}}}$ comes from the factorization of the amplitude of $B^- \to X_{c\bar{c}} K^-$, followed by the strong decay $X_{c\bar{c}} \to \Lambda_c^+ \bar{\Lambda}_c^-$. Since the chirality-flipping baryon vector form factor $f_2(q^2)$ is in general suppressed by two more powers of the dibaryon invariant mass q^2 than $f_1(q^2)$ and since $q^2 \sim 4m_{\Lambda_c}^2$ is large in $B \to \Lambda_c \bar{\Lambda}_c K$ decays, we expect the contributions from $X_{c\bar{c}}$ coupled to $\Lambda_c \bar{\Lambda}_c$ through $h_2^{\Lambda_c \bar{\Lambda}_c V}$ be small and hence can be neglected in our calculation. The $h_2^{\Lambda_c \bar{\Lambda}_c A}$ term in the axial-vector decay amplitude can also be dropped since $q_\nu(-g^{\mu\nu} + q^\mu q^\nu/m_A^2) \sim 0$ due to the fact that $m_A^2 \sim q^2$ within the phase space. The decay amplitude $A_{X_{c\bar{c}}}$ then becomes

$$A_{V(A)}(B^- \to \Lambda_c^+ \bar{\Lambda}_c^- K^-) = \frac{G_F}{\sqrt{2}} V_{cs} V_{cb}^* a_2 m_{V(A)} f_{V(A)} \bar{u}_{\Lambda_c}$$
$$\times [\mathcal{M}_{V(A)} \slashed{p}_K (\gamma_5)$$
$$+ \mathcal{M}_{S(P)} (\gamma_5)] v_{\bar{\Lambda}_c}, \quad (14)$$

with

$$\mathcal{M}_V = 2F_1^{BK}(q^2) \left(\frac{-h_1^{\Lambda_c \bar{\Lambda}_c V}}{q^2 - m_V^2 + i m_V \Gamma_V} \right), \qquad \mathcal{M}_S = 0,$$

$$\mathcal{M}_A = 2F_1^{BK}(q^2) \left(\frac{-h_1^{\Lambda_c \bar{\Lambda}_c A}}{q^2 - m_A^2 + i m_A \Gamma_A} \right),$$

$$\mathcal{M}_P = \left[F_1^{BK}(q^2) \frac{q^2 - (m_B^2 - m_K^2)}{q^2} + F_0^{BK}(q^2)(m_B^2 - m_K^2) \right.$$
$$\left. \times \left(\frac{1}{q^2} - \frac{1}{m_A^2} \right) \right] 2m_{\Lambda_c} \left(\frac{-h_1^{\Lambda_c \bar{\Lambda}_c A}}{q^2 - m_A^2 + i m_A \Gamma_A} \right). \quad (15)$$

The minus signs in front of the strong couplings $h_1^{\Lambda_c \bar{\Lambda}_c V,A}$ come from the minus sign of the $g^{\mu\nu}$ part of the $X_{c\bar{c}}$ propagator in Eq. (11).

We take $f_V \cdot h_1^{\Lambda_c \bar{\Lambda}_c V} = f_A \cdot h_1^{\Lambda_c \bar{\Lambda}_c A} = 4$ GeV and show in Fig. 2 the plots of branching ratios for each kind of resonance as functions of both the mass and the width of the resonance. The branching fractions depend on $(f_{X_{c\bar{c}}} \cdot h_1^{\Lambda_c \bar{\Lambda}_c V})^2$ and $(f_{X_{c\bar{c}}} \cdot h_1^{\Lambda_c \bar{\Lambda}_c A})^2$ proportionally. We notice that the width of the axial-vector resonance Γ_A is more than an order of magnitude smaller than Γ_V when both $\mathcal{B}_V$ and $\mathcal{B}_A$ are around the experimental value $\mathcal{B} \sim 7 \times 10^{-4}$. This is due mainly to the smallness of $\mathcal{B}_A$ which suffers from the destructive interference between *comparable* $\mathcal{M}_A$ and $\mathcal{M}_P$ contributions in the decay rate. Figure 3 shows the decay rates from $\mathcal{M}_A$, $\mathcal{M}_P$, and from $\text{Re}\mathcal{M}_A \mathcal{M}_P^*$ alone without taking into account the resonance effect. One can see that the $\mathcal{M}_A$ and $\mathcal{M}_P$ contributions are comparable while the interference term $\text{Re}\mathcal{M}_A \mathcal{M}_P^*$ gives almost twice the negative contribution of either $\mathcal{M}_A$ or $\mathcal{M}_P$, resulting in a large cancellation with $\mathcal{M}_A$ and $\mathcal{M}_P$. Therefore, in order to counteract this cancellation, one needs a smaller width of the resonance such that $|1/(q^2 - m_A^2 + i m_A \Gamma_A)|^{-2}$ becomes more singular in the allowed range of q^2. There is, however, no such interference found in $\mathcal{B}_V$ as $\mathcal{M}_V$ stands
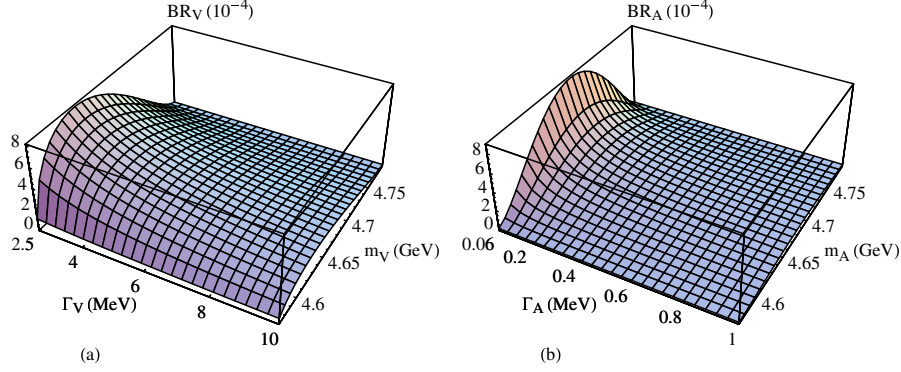
FIG. 2 (color online).   Branching fractions of $\mathcal{B}_{X_{c\bar{c}}}(B^- \to \Lambda_c^+ \bar{\Lambda}_c^- K^-)$ as a function of the mass and the decay width of the intermediate resonance $X_{c\bar{c}}$ when it is (a) a vector ($X_{c\bar{c}} = V$) and (b) an axial-vector ($X_{c\bar{c}} = A$) particle with $f_V \cdot h_1^{\Lambda_c \bar{\Lambda}_c V} = f_A \cdot h_1^{\Lambda_c \bar{\Lambda}_c A} = 4$ GeV.

alone in the vector-induced decay amplitude after taking $h_2^{\Lambda_c \bar{\Lambda}_c V} = 0$ in Eq. (12).

Therefore, the above analysis seems to imply the existence of a narrow hidden charm bound state with a mass of order $4.6 \sim 4.7$ GeV that couples strongly with the charmed baryon pair. Recall that many new charmonium-like resonances with masses around 4 GeV starting with $X(3872)$ [13] and so far ending with $Y(4260)$ [14] have been recently observed by *BABAR* and Belle. These charmoniumlike states are above the $D\bar{D}$ threshold but below the two-charmed baryon threshold. The new state we have put forward is just marginally above the $\Lambda_c \bar{\Lambda}_c$ threshold. In principle, this new state can be searched for in B decays and in $p\bar{p}$ collisions by studying the mass spectrum of $D^{(*)} \bar{D}^{(*)}$ or $\Lambda_c \bar{\Lambda}_c$.

III. TWO-BODY DECAYS

The two-body doubly charmed baryonic B decays $B^- \to \Xi_c^0 \bar{\Lambda}_c^-$ and $\bar{B}^0 \to \Xi_c^+ \bar{\Lambda}_c^-$ receive contributions from the internal W emission (see Fig. 4) and weak annihilation. The latter contribution can be safely neglected as it is not only quark mixing but also helicity suppressed. It should be stressed that, in contrast to the internal W emission in mesonic B decays, internal W emission in baryonic B decay is not necessarily color suppressed. This is because the baryon wave function is totally antisymmetric in color indices. One can see from Fig. 4 that there is no color suppression for the meson production. In the effective Hamiltonian approach, the relevant weak Hamiltonian is

$$\mathcal{H}_{\text{eff}} = \frac{G_F}{\sqrt{2}} V_{cb} V_{us}^* (c_1 O_1 + c_2 O_2)$$

$$\to \frac{G_F}{\sqrt{2}} V_{cb} V_{us}^* (c_1 - c_2) O_1, \tag{16}$$

where $O_1 = (\bar{c}b)(\bar{s}c)$ and $O_2 = (\bar{c}c)(\bar{s}b)$. In the above equation, we have used the fact that the operator $O_1 - O_2$ is antisymmetric in color indices (more precisely, it is a color antitriplet). Therefore, the Wilson coefficient for the tree-dominated internal W emission is $c_1 - c_2$ rather than $a_2 = c_2 + c_1/3$. This is indeed the case found in the pole model calculation in [15].
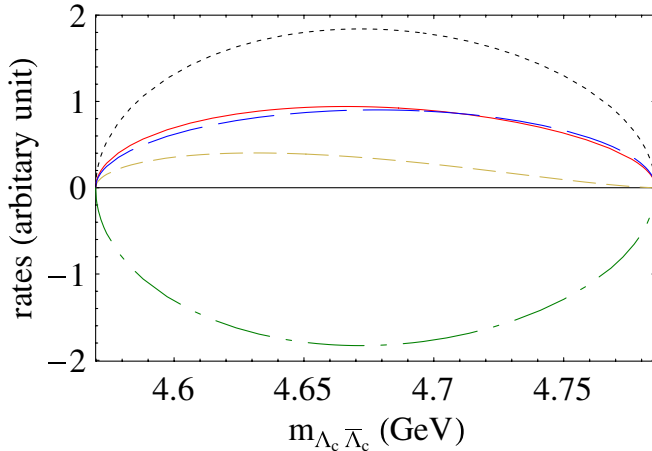

FIG. 3 (color online).   Decay rates contributed from $\mathcal{M}_V$ (short-dashed line), $\mathcal{M}_A$ (solid line), $\mathcal{M}_P$ (long-dashed line), sum of the previous two (dotted line), and from $2\,\text{Re}\,\mathcal{M}_A \mathcal{M}_P^*$ (dot-dashed line) alone without taking into account the resonance effect, i.e. $(q^2 - m_A^2 + i m_A \Gamma_A)^{-1}$ is replaced by any constant in each term.
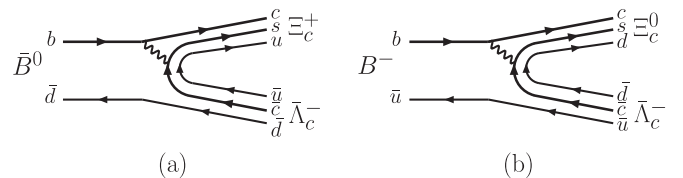

FIG. 4.   (a) $\bar{B}^0 \to \Xi_c^+ \bar{\Lambda}_c^-$ and (b) $B^- \to \Xi_c^0 \bar{\Lambda}_c^-$ as proceeding through internal W-emission diagrams.

Since the internal *W* emission in Fig. 4 is not factorizable, it is difficult to evaluate its amplitude directly. Pole model has been applied in [15] to compute $B \to \mathcal{B}_1 \bar{\mathcal{B}}_2$. However, the strong coupling involved in this model is unknown and hence it has to be fixed from other processes, e.g. the 3-body baryonic *B* decays. Since the CKM angles for $\bar{B}^0 \to \Xi_c^+ \bar{\Lambda}_c^-$ and $\bar{B}^0 \to \Lambda_c^+ \bar{p}$ are the same in magnitude (but opposite in sign), the pole model does not explicitly explain why the former has a rate much larger than the latter. In particular, the dynamical suppression of $\Lambda_c^+ \bar{p}$ relative to $\Xi_c^+ \bar{\Lambda}_c^-$ is not clearly manifested in the pole model calculation. In order to understand why $\Xi_c^+ \bar{\Lambda}_c^- \gg \Lambda_c^+ \bar{p}$, let us reexamine Fig. 4.

There are several possibilities for the quark-antiquark pair creation in Fig. 4. In one case, $q\bar{q}$ is picked up from the vacuum via the soft nonperturbative interactions so that it carries the vacuum quantum numbers $^3P_0$. It is also possible that the quark pair is created perturbatively via one-gluon exchange with one-gluon quantum numbers $^3S_1$. It is not clear which mechanism, the $^3P_0$ or $^3S_1$ model, dominates the 2-body baryonic *B* decays, though in practice the $^3P_0$ model is simpler. Since the energy release is relatively small in charmful baryonic *B* decay, the $^3P_0$ model seems to be more relevant. In the present work, we also consider the possibility that the $q\bar{q}$ pair is produced via a light meson exchange. The $q\bar{q}$ pair created from soft nonperturbative interactions tends to be soft. For an energetic proton produced in 2-body *B* decays, the momentum fraction carried by its quark is large, $\sim \mathcal{O}(1)$, while for an energetic charmed baryon, its momentum is carried mostly by the charmed quark. As a consequence, the doubly charmed baryon state such as $\Xi_c \bar{\Lambda}_c$ has a configuration more favorable than $\Lambda_c \bar{p}$.

In order to evaluate Fig. 4 for the decay $\bar{B} \to \Xi_c \bar{\Lambda}_c$, we need to know the distribution amplitudes of the charmed baryon $\mathcal{B}_c$ and the *B* meson. For the wave functions of $\mathcal{B}_c = \Xi_c^0, \Lambda_c$, they have the forms [16,17]

$$\langle \Xi_c^0(p) | \bar{c}_\alpha^a(z_1) \bar{s}_\beta^b(z_2) \bar{d}_\gamma^c(z_3) | 0 \rangle$$

$$= \frac{\epsilon^{abc}}{6} \frac{f_{\Xi_c}}{4} [\bar{u}_{\Xi_c^0}(p)]_\alpha [C^{-1}\gamma_5(\not{p} + m_{\Xi_c})]_{\gamma\beta} \Psi_{\Xi_c^0}(z_1, z_2, z_3),$$

$$\langle \bar{\Lambda}_c(p') | c_\alpha^a(z_1') u_\beta^b(z_2') d_\gamma^c(z_3') | 0 \rangle$$

$$= \frac{\epsilon^{abc}}{6} \frac{f_{\Lambda_c}}{4} [\bar{v}_{\Lambda_c}(p')]_\alpha [(\not{p}' - m_{\Lambda_c})\gamma_5 C]_{\beta\gamma} \Psi_{\Lambda_c}(z_1', z_2', z_3'),$$

$$(17)$$

where *c*, *q*, and *d* are the quark fields, *a*, *b*, and *c* the color indices, α, β, and γ the spinor indices, *C* the charge conjugation matrix, and $f_{\mathcal{B}_c}$ the decay constant. Following [18] we can write

$$\Psi_{\Xi_c}(z_1, z_2, z_3) = \int [dx][d^2 k_\perp]$$

$$\times e^{i\Sigma k_i \cdot z_i} \Psi_{\Xi_c}(x_1, x_2, x_3, \mathbf{k}_{1\perp}, \mathbf{k}_{2\perp}, \mathbf{k}_{3\perp}),$$

$$\Psi_{\Lambda_c}(z_1', z_2', z_3') = \int [dx'][d^2 k'_\perp]$$

$$\times e^{i\Sigma k'_i \cdot z_i} \Psi_{\Lambda_c}(x_1', x_2', x_3', \mathbf{k}'_{1\perp}, \mathbf{k}'_{2\perp}, \mathbf{k}'_{3\perp}),$$

$$(18)$$

where $p^{(\prime)} = (p^{(\prime)+}, p^{(\prime)-}, 0_\perp)$ is the momentum of $\Xi_c^0(\bar{\Lambda}_c^-)$ and $k_1^{(\prime)}, k_2^{(\prime)}, k_3^{(\prime)}$ the momenta of the constituent quarks of the baryons, which are taken to be [17,19]

$$k_1 = (x_1 p^+, p^-, \mathbf{k}_{1\perp}), \quad k_1' = (p'^+, x_1' p'^-, \mathbf{k}'_{1\perp}),$$

$$k_{2(3)} = (x_{2(3)} p^+, 0, \mathbf{k}_{2(3)\perp}), \quad k'_{2(3)} = (0, x'_{2(3)} p'^-, \mathbf{k}'_{2(3)\perp})$$

$$(19)$$

$$[dx^{(\prime)}] = dx_1^{(\prime)} dx_2^{(\prime)} dx_3^{(\prime)} \delta\left(1 - \sum_{i=1}^3 x_i^{(\prime)}\right),$$

$$[d^2 k_\perp^{(\prime)}] = d^2 k_{1\perp}^{(\prime)} d^2 k_{2\perp}^{(\prime)} d^2 k_{3\perp}^{(\prime)} \delta^2(\mathbf{k}_{1\perp}^{(\prime)} + \mathbf{k}_{2\perp}^{(\prime)} + \mathbf{k}_{3\perp}^{(\prime)}),$$

$$(20)$$

with x_i being the momentum fractions associated with the baryon, and $\mathbf{k}_{i\perp}$ the corresponding transverse momenta. Note that for simplicity, light quark masses are neglected in Eq. (19). The *B* meson wave function is expressed as

$$\langle 0 | \bar{q}_\alpha^b(z_1) b_\beta^a(z_2) | B(p) \rangle = -i \frac{\delta^{ab}}{3} \frac{f_B}{4} [(\not{p} + m_B)\gamma_5]_{\beta\alpha}$$

$$\times \int_0^1 d\xi e^{-i\Sigma p_i \cdot z_i} \Phi_B(\xi),$$

$$(21)$$

with $p_2 = p_b = ((1-\xi)p_B^+, (1-\xi)p_B^-, \mathbf{0}_\perp)$, $p_1 = p_l = (\xi p_B^+, \xi p_B^-, \mathbf{0}_\perp)$ and $p_B^+ = p_B^- = m_B$.

The $B^- \to \Xi_c^0 \bar{\Lambda}_c^-$ decay amplitude now consists of three parts corresponding to the exchange of the σ, π^0, and π^- between the soft $q\bar{q}$ quark pair and the spectator as shown in Fig. 5,

$$A(B^- \to \Xi_c^0 \bar{\Lambda}_c^-) = A_\sigma + A_{\pi^0} + A_{\pi^-},$$

where, for instance, in the case that the decay proceeds through the exchange of σ or π^0

$$A_{\sigma(\pi^0)} = \frac{G_F}{\sqrt{2}} V_{cb} V_{us}^* (c_1 - c_2)$$

$$\times \int dz dz' e^{i(k_3 + k_3') \cdot z} e^{i(k_2' - p_\ell) \cdot z'} g_{\sigma(\pi)qq}^2 D_F(z - z')$$

$$\times \langle \Xi_c^0 | \bar{c}_\alpha^a(0) \bar{s}_\beta^b(0) \bar{d}_\gamma^c(z) | 0 \rangle$$

$$\times \langle \bar{\Lambda}_c | c_\delta^b(0) u_{\eta'}^d(z') d_{\gamma'}^c(z) | 0 \rangle \langle 0 | \bar{u}_\eta^d(z') b_\rho^a(0) | B^- \rangle$$

$$\times [\gamma_\mu (1 - \gamma_5)]_{\alpha\rho} [\gamma^\mu (1 - \gamma_5)]_{\beta\delta} \Gamma_{\gamma\gamma'} \Gamma_{\eta\eta'}$$
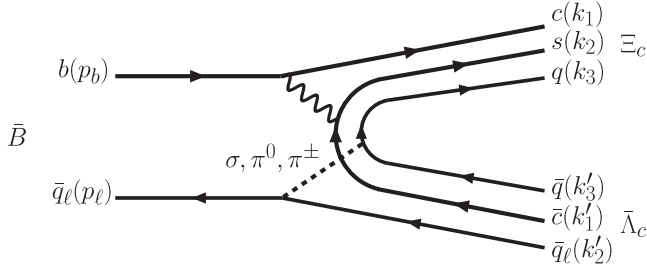
$$(22)$$

FIG. 5. Soft $q_{(\ell)}\bar{q}_{(\ell)}[= u\bar{u}(d\bar{d}), d\bar{d}(u\bar{u})$ and $d\bar{u}(u\bar{d})]$ as produced through the σ, π^0, and $\pi^\pm$ meson exchanges. Inside the parentheses are the momenta that the constituent quarks are carrying. Note the arrows on the quark lines do not represent the momentum flows.

$$
\begin{aligned}
= &\frac{G_F}{\sqrt{2}}V_{cb}V_{us}^* \frac{f_B f_{\Xi_c} f_{\Lambda_c}}{64}\frac{(c_1 - c_2)}{18}\int d\xi \int [dx][dx'][d^2k_\perp] \\
&\times [d^2k_\perp'](2\pi)^4\delta^4(p_\ell - k_2' - k_3' - k_3) \\
&\times \Psi_{\Xi_c}(x, \mathbf{k}_\perp)\Psi_{\Lambda_c}(x', \mathbf{k}_\perp')\Phi_B(\xi) \\
&\times \frac{g^2_{\sigma(\pi)qq}}{p^2_{\sigma(\pi)} - m^2_{\sigma(\pi)} + im_{\sigma(\pi)}\Gamma_{\sigma(\pi)}}\bar{u}_{\Xi_c}[a_{\sigma(\pi)} + b_{\sigma(\pi)}\gamma_5]v_{\Lambda_c}
\end{aligned}
$$
(23)

with $\Gamma_{\eta\eta'} = \Gamma_{\gamma\gamma'} = 1$ for σ and $\Gamma_{\eta\eta'} = -\Gamma_{\gamma\gamma'} = i\gamma_5$ for π^0 in Eq. (22). Note that the factor of $1/18$ in Eq. (23) is the color factor. For the $\pi^\pm$ exchange, the terms $g_{\pi qq}$, $\Gamma_{\gamma\gamma'}$, and $\langle\bar{\Lambda}_c|c^b_\delta(0)u^d_{\eta'}(z')d^c_{\gamma'}(z)|0\rangle$ in A_{π^0} are replaced by $\sqrt{2}g_{\pi qq}$, $-\Gamma_{\gamma\gamma'}$, and $\langle\bar{\Lambda}_c|c^b_\delta(0)d^d_{\eta'}(z')u^c_{\gamma'}(z)|0\rangle = -\langle\bar{\Lambda}_c|c^b_\delta(0)u^d_{\eta'}(z)d^c_{\gamma'}(z')|0\rangle$, respectively, where Eq. (17) has been used. Because of the symmetry property of the Λ_c wave function given in Eq. (28), the $\pi^\pm$ contribution is the same as the π^0 one except for an enhancement factor of $(\sqrt{2})^2$ arising from isospin. Because of the tiny mass difference between π^0 and π^- and the extremely narrow widths of these two particles, we have $A_{\pi^-}/A_{\pi^0} = 2$ to a very good precision. The momentum labels of the quarks in Ξ_c and $\bar{\Lambda}_c$ are depicted in Fig. 5, p_b and p_ℓ are the momenta of the b quark and the light spectator quark of the B meson defined after Eq. (21), respectively, $g_{\sigma(\pi)qq}$ is the coupling of the $\sigma(\pi)$ meson with the $q\bar{q}$ pair, and

$$
\begin{aligned}
a_\sigma &= -4m_{\Xi_c}[(m_B + m_{\Lambda_c})^2 - m^2_{\Xi_c}], \\
b_\sigma &= -4m_{\Xi_c}[(m_{\Lambda_c} + m_{\Xi_c})^2 - m^2_B],
\end{aligned}
$$
(24)

for σ and

$$
\begin{aligned}
a_\pi &= 4m_{\Xi_c}[m^2_B - (m_{\Xi_c} - m_{\Lambda_c})^2], \\
b_\pi &= 4m_{\Xi_c}[m^2_{\Xi_c} - (m_B - m_{\Lambda_c})^2],
\end{aligned}
$$
(25)

for both π^0 and π^-.

Our results are consistent with heavy quark effective theory. It has been shown that in the heavy quark limit,

the decay amplitude can be expressed as $\bar{u}[A + B\gamma_5]v$ with [20]

$$
\begin{aligned}
A &= 2\sqrt{m_B}[\alpha(r_1 - r_2) - \beta], \\
B &= -2\sqrt{m_B}[\alpha(r_1 + r_2) - \beta],
\end{aligned}
$$
(26)

where $r_1 = m_{\Xi_c}/m_B$, $r_2 = m_{\Lambda_c}/m_B$ and α, β are two unknown parameters. It is easily seen that, Eqs. (24)–(26) agree with each other after setting $\alpha = \alpha_\sigma + \alpha_\pi$, $\beta = \beta_\sigma + \beta_\pi$, with $\alpha_\sigma = \beta_\sigma \propto 1 + r_1 + r_2$ for σ and $\alpha_\pi = \beta_\pi \propto 1 + r_1 - r_2$ for both π^0 and π^-, where the overall coefficients can be easily determined from Eqs. (23)–(25).

To proceed with the numerical calculations, we need to first deal with the delta functions that impose constraints on the integration limits as well as relations between integral variables. We show in the appendix the decay amplitude and the integral variables as a result of the delta function integrations. The wave functions are adopted to be

$$
\begin{aligned}
\Psi_{\mathcal{B}_c}(x_1, x_2, x_3) &= \int[d^2k_\perp]\Psi_{\mathcal{B}_c}(x_i, \mathbf{k}_{i\perp}) \\
&= N_{\mathcal{B}_c}x_1x_2x_3\exp\Big[-\frac{\hat{m}_c^2}{2\beta^2 x_1} - \frac{\hat{m}_2^2}{2\beta^2 x_2} \\
&\quad - \frac{\hat{m}_3^2}{2\beta^2 x_3}\Big],
\end{aligned}
$$
(27)

with

$$
\begin{aligned}
\Psi_{\mathcal{B}_c}(x_i, \mathbf{k}_{i\perp}) &= \frac{N_{\mathcal{B}_c}}{(2\pi\beta^2)^2}\exp\Big[-\frac{\mathbf{k}^2_{1\perp} + \hat{m}^2_c}{2\beta^2 x_1} - \frac{\mathbf{k}^2_{2\perp} + \hat{m}^2_2}{2\beta^2 x_2} \\
&\quad - \frac{\mathbf{k}^2_{3\perp} + \hat{m}^2_3}{2\beta^2 x_3}\Big],
\end{aligned}
$$
(28)

for the charmed baryon [21] and

$$
\Phi_B(x) = N_B x^2(1 - x)^2\exp\Big[-\frac{1}{2}\frac{x^2 m^2_B}{\omega^2_B}\Big]
$$
(29)

for the B meson [22]. The wave functions obey the normalization

$$
\int[dx]\Psi(x_1, x_2, x_3) = 1, \qquad \int_0^1 dx\Phi_B(x) = 1.
$$
(30)

The other input parameters are specified as follows. The decay constants for the charmed baryons can be related to that of the Λ_b baryon via the relation [19]

$$
f_{\mathcal{B}_c}m_{\mathcal{B}_c} = f_{\Lambda_b}m_{\Lambda_b},
$$
(31)

valid in the heavy quark limit. Using $f_{\Lambda_b} = 2.71 \times 10^{-3}$ GeV^2 obtained from a fit of the perturbative quantum chromodynamics (PQCD) calculation for $\Lambda_b \to \Lambda_c$ decays to $\mathcal{B}(\Lambda_b \to \Lambda_c l\bar{\nu})$ [19], it is found that $f_{\Lambda_c} = 6.7 \times 10^{-3}$ GeV^2 and $f_{\Xi_c} = 6.2 \times 10^{-3}$ GeV^2, which are roughly ~ 1.3–2.3 times that of the results from QCD

TABLE I. Predictions on the branching ratios of $B^- \to \Xi_c^0 \bar{\Lambda}_c^-$ and $\bar{B}^0 \to \Xi_c^+ \bar{\Lambda}_c^-$ decays. The first and second errors come from the theoretical uncertainties in the parameters β and ω_b, respectively, which are taken to be $\beta = 1.20 \pm 0.05$ GeV and $\omega_b = 0.40 \pm 0.05$ GeV, and the third error from the baryon decay constants. Results shown in second and third raws are from π or σ exchange alone, respectively.

Mode	Theory (10^{-3})	Expt (10^{-3})	Mode	Theory (10^{-3})	Expt (10^{-3})
$\mathcal{B}(B^- \to \Xi_c^0 \bar{\Lambda}_c^-)$	$2.2^{+0.6+5.1+6.1}_{-0.6-1.9-1.9}$	≈ 4.8	$\mathcal{B}(\bar{B}^0 \to \Xi_c^+ \bar{\Lambda}_c^-)$	$2.0^{+0.5+4.7+5.6}_{-0.6-1.7-1.7}$	≈ 1.2
$\mathcal{B}(B^- \to \Xi_c^0 \bar{\Lambda}_c^-)_\pi$	$1.8^{+0.5+4.2+5.0}_{-0.5-1.6-1.6}$		$\mathcal{B}(\bar{B}^0 \to \Xi_c^+ \bar{\Lambda}_c^-)_\pi$	$1.7^{+0.5+3.9+4.7}_{-0.5-1.4-1.5}$	
$\mathcal{B}(B^- \to \Xi_c^0 \bar{\Lambda}_c^-)_\sigma$	$0.2^{+0.0+0.4+0.6}_{-0.1-0.1-0.2}$		$\mathcal{B}(\bar{B}^0 \to \Xi_c^+ \bar{\Lambda}_c^-)_\sigma$	$0.2^{+0.0+0.4+0.6}_{-0.0-0.1-0.2}$	

sum rules [23].[2] In our calculations we shall employ the value of the Λ_b baryon decay constant which is in the middle of the range that has the PQCD value as the lower bound and the highest value from QCD sum rules as the upper bound. The deviations of the lower and the upper bounds from this central value are then taken as one of the theoretical errors in our model.

For the *B* meson, we use $f_B = 0.2$ GeV. For the coupling $g_{\sigma(\pi)qq}$, the linear sigma model leads to $g_{\sigma NN} = \sqrt{2} m_N / f_\pi$ with $f_\pi = 132$ MeV, and $g_{\pi NN} = \sqrt{2} m_N g_A / f_\pi$ with $g_A \simeq 1.25$ from the Goldberger-Treiman relation. Hence, it is reasonable to take $g_{\sigma(\pi)qq} = g_{\sigma(\pi)NN}/3$ in the constituent quark model. For the σ meson, we use $\Gamma_\sigma \approx m_\sigma = 600$ MeV [4]. The constituent quark masses appearing in the $\mathcal{B}_c$ wave function are taken to be $\hat{m}_u = \hat{m}_d = 0.33$ GeV and $\hat{m}_s = 0.55$ GeV [21], while $\hat{m}_c = m_{\mathcal{B}_c}$ is employed.

The decay rate is given by

$$\Gamma(B \to \mathcal{B}_1 \bar{\mathcal{B}}_2) = \frac{p_c}{4\pi} \left\{ |A|^2 \frac{(m_B + m_1 + m_2)^2 p_c^2}{(E_1 + m_1)(E_2 + m_2) m_B^2} \right.$$
$$\left. + |B|^2 \frac{[(E_1 + m_1)(E_2 + m_2) + p_c^2]^2}{(E_1 + m_1)(E_2 + m_2) m_B^2} \right\},$$
(32)

where p_c is the c.m. momentum, E_i and m_i are the energy and mass of the baryon $\mathcal{B}_i$, respectively. The results of calculations are summarized in Table I. The theoretical errors come from the uncertainties in the parameters β, ω_b, which are taken to be $\beta = 1.20 \pm 0.05$ GeV,[3] $\omega_b = 0.40 \pm 0.05$ GeV, and the baryon decay constants. It is clear that the pion exchange gives the dominant contribution owing to its narrow width. The prediction is in agreement with experiment for $\Xi_c^+ \bar{\Lambda}_c^-$, but a bit small for $\Xi_c^0 \bar{\Lambda}_c^-$

------

[2]The Λ_b decay constant is found to be in the range $(2.0–3.5) \times 10^{-2}$ GeV^3 in QCD sum rules [23]. It differs from the Λ_b decay constant in this work by a factor of Λ_b mass. After normalizing it to having the same dimension as the decay constants in this work, the above range turns out to be about 1.3–2.3 times that of the decay constant from PQCD.

[3]As shown in [21], the parameter β is of order 1 GeV for light baryons. Just as the meson case [11], β should become larger for the heavy baryons.

The above calculation is not applicable to the two-body decay $\bar{B}^0 \to \Lambda_c^+ \bar{p}$ with one charmed baryon in the final state. This is because two hard gluons are needed to produce an energetic antiproton: one hard gluon for kicking the spectator quark of the *B* meson to make it energetic and the other for producing the hard $q\bar{q}$ pair. The pQCD calculation for this decay will be much more involved (see e.g. [17] for pQCD calculations of $\Lambda_b \to \Lambda J/\psi$) and is beyond the scope of the present work. Nevertheless, it is expected that $\Gamma(\bar{B} \to \mathcal{B}_c \bar{N}) \ll \Gamma(\bar{B} \to \Xi_c \bar{\Lambda}_c)$ as the former is suppressed by order of α_s^4. This dynamical suppression effect for the $\Lambda_c \bar{p}$ production relative to $\Xi_c \bar{\Lambda}_c$ has been neglected in the previous studies based on QCD sum rules [9] and on the diquark model [10].

IV. CONCLUSIONS

In this work we have studied the two-body and three-body doubly charmed baryonic *B* decays, namely, $\bar{B} \to \Xi_c \bar{\Lambda}_c$ and $\bar{B} \to \Lambda_c \bar{\Lambda}_c K$, aiming to resolve the puzzles associated with them. We point out that the suppression from the $\Lambda_c^+ \bar{\Lambda}_c^- K^-$ phase space is so strong that $\Lambda_c^+ \bar{\Lambda}_c^-$ pair is unlikely to be produced dominantly through the direct three-body decay processes. Nevertheless, the great suppression for the $\Lambda_c^+ \bar{\Lambda}_c^- K$ production can be alleviated provided that there exists a narrow hidden charm bound state with a mass near the $\Lambda_c \bar{\Lambda}_c$ threshold, of order $4.6 \sim 4.7$ GeV. This new state that couples strongly to the charmed baryon pair can be searched for in *B* decays and in $p\bar{p}$ collisions by studying the mass spectrum of $D^{(*)} \bar{D}^{(*)}$ or $\Lambda_c \bar{\Lambda}_c$.

The doubly charmful decay such as $\bar{B} \to \Xi_c \bar{\Lambda}_c$ has a configuration more favorable than the singly charmful one such as $\bar{B}^0 \to \Lambda_c \bar{p}$ even though they have the same CKM angles in magnitude. This is because no hard gluon is needed to produce the energetic $\Xi_c \bar{\Lambda}_c$ pair in the former decay, while two hard gluons are needed for the latter process. Therefore, $\Lambda_c \bar{p}$ is suppressed relative to $\Xi_c \bar{\Lambda}_c$ due to a dynamical suppression from $\mathcal{O}(\alpha_s^4)$. Assuming that a soft $q\bar{q}$ quark pair is produced through the σ and π meson exchanges in the configuration for $\bar{B} \to \Xi_c \bar{\Lambda}_c$, it is found that $\mathcal{B}(\bar{B} \to \Xi_c \bar{\Lambda}_c) \sim 10^{-3}$.

Note added.—After this paper was submitted for publication, Belle published the updated version of [1], in which the spectrum of the $\Lambda_c\bar{\Lambda}_c$ pair in the $\bar{B} \to \Lambda_c\bar{\Lambda}_c K$ decays is shown for the first time. According to Belle's observation no new resonance with a mass near the $\Lambda_c\bar{\Lambda}_c$ threshold was found (see Fig. 3 in version 2 of [1]). This implies the failure of naive factorization for this decay mode and may hint at the importance of nonfactorizable contributions such as final-state effects. For example, the weak decay $B \to D^{(*)}\bar{D}_s^{(*)}$ followed by the rescattering $D^{(*)}\bar{D}_s^{(*)} \to \Lambda_c\bar{\Lambda}_c K$ [24] or the decay $B \to \Xi_c\bar{\Lambda}_c$ followed by $\Xi_c\bar{\Lambda}_c \to \Lambda_c\bar{\Lambda}_c K$ may explain the large rate observed for $B \to \Lambda_c\bar{\Lambda}_c K$.

APPENDIX

After integrating out the delta functions in Eq. (23), we obtain, taking decay via σ as an example,

$$
\begin{aligned}
A_\sigma(B^- \to \Xi_c^0\bar{\Lambda}_c^-) = {} & \frac{G_F}{\sqrt{2}} V_{cb}V_{us}^* \frac{f_B f_{\Xi_c} f_{\Lambda_c}}{1152}(c_1 - c_2)(2\pi)^4 2 \int_0^{p'^-/p_B^-} d\xi \int_0^{1-\xi p_B^+/p^+} \frac{dx_2}{p^+} \int_0^{\xi p_B^-/p'^-} \frac{dx_2'}{p'^-}(2\pi)\left(\frac{1}{2}\right)^3 \\
& \times \int_0^\infty d\mathbf{k}_{3\perp}^2 \int_0^\infty d\mathbf{k}_{2\perp}^2 \int_0^{2\pi} d\theta_{23} \int_0^\infty d\mathbf{k}_{3\perp}'^2 \int_0^{2\pi} d\theta_{33'} \Psi_{\Xi_c}(x, \mathbf{k}_\perp)\Psi_{\Lambda_c}(x', \mathbf{k}_\perp')\Phi_B(\xi) \\
& \times \frac{(g_{\sigma qq})^2}{p_\sigma^2 - m_\sigma^2 + im_\sigma\Gamma_\sigma} \bar{u}_{\Xi_c}(a_\sigma + b_\sigma\gamma_5)v_{\Lambda_c},
\end{aligned}
\tag{A1}
$$

where $p^{(\prime)} = (p^{(\prime)+}, p^{(\prime)-}, 0_\perp)$ is the momentum of $\Xi_c^0(\bar{\Lambda}_c^-)$ and $k_1^{(\prime)}, k_2^{(\prime)}, k_3^{(\prime)}$ the momenta of the constituent quarks of the baryons, which are taken to be

$$
\begin{aligned}
k_1 &= (x_1 p^+, p^-, \mathbf{k}_{1\perp}), \quad k_1' = (p'^+, x_1' p'^-, \mathbf{k}_{1\perp}'), \\
k_{2(3)} &= (x_{2(3)} p^+, 0, \mathbf{k}_{2(3)\perp}), \quad k_{2(3)}' = (0, x_{2(3)}' p'^-, \mathbf{k}_{2(3)\perp}').
\end{aligned}
\tag{A2}
$$

The integrations of the delta functions yield the following relations

$$
\begin{aligned}
x_1^{(\prime)} &= 1 - x_2^{(\prime)} - x_3^{(\prime)}, \qquad x_3 = \frac{\xi p_B^+}{p^+}, \\
x_3' &= \frac{\xi p_B^-}{p'^-} - x_2', \qquad \mathbf{k}_{1\perp}^{(\prime)} = -(\mathbf{k}_{2\perp}^{(\prime)} + \mathbf{k}_{3\perp}^{(\prime)}), \\
\mathbf{k}_{2\perp}' &= -(\mathbf{k}_{3\perp} + \mathbf{k}_{1\perp}') \Rightarrow \mathbf{k}_{1\perp}' = \mathbf{k}_{3\perp},
\end{aligned}
\tag{A3}
$$

and the limits of integrations as shown in Eq. (A1). Thus,

$$
\begin{aligned}
p_\sigma^2 &= (k_3 + k_3')^2 \\
&= \left(\frac{\xi p_B^+}{p^+}\right)\left(\frac{\xi p_B^-}{p'^-} - x_2'\right)p^+ p^- - (\mathbf{k}_{3\perp}^2 + \mathbf{k}_{3\perp}'^2 \\
&\quad + 2|\mathbf{k}_{3\perp}||\mathbf{k}_{3\perp}'|\cos\theta_{33'}), \\
\mathbf{k}_{1\perp}^2 &= \mathbf{k}_{2\perp}^2 + \mathbf{k}_{3\perp}^2 + 2|\mathbf{k}_{2\perp}||\mathbf{k}_{3\perp}|\cos\theta_{23}, \\
\mathbf{k}_{2\perp}'^2 &= \mathbf{k}_{3\perp}^2 + \mathbf{k}_{3\perp}'^2 + 2|\mathbf{k}_{3\perp}||\mathbf{k}_{3\perp}'|\cos\theta_{33'},
\end{aligned}
\tag{A4}
$$

where θ_{23} and $\theta_{33'}$ are the angles of $\mathbf{k}_{2\perp}$ and of $\mathbf{k}_{3\perp}'$ as measured against $\mathbf{k}_{3\perp}$, respectively. Note that the integration ranges of x_2 and x_2' are constrained by the delta function in Eq. (23), while the integration range of ξ is restricted by $(1 - \xi p_B^+/p^+) \geq 0$ and $\xi p_B^-/p'^- \leq 1$.

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