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Charmful three-body baryonic B decays

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We study the charmful three-body baryonic B decays with $D^{(*)}$ or J/Ψ in the final state. We explain the measured rates of $\bar{B}^0 \to n\bar{p}D^{*+}$, $\bar{B}^0 \to p\bar{p}D^{(*)0}$, and $B^- \to \Lambda\bar{p}J/\Psi$ and predict the branching fractions of $\bar{B}^0 \to \Lambda\bar{p}D^{*+}$, $\bar{B}^0 \to \Sigma^0\bar{p}D^{*+}$, $B^- \to \Lambda\bar{p}D^0$, and $B^- \to \Lambda\bar{p}D^{*0}$ to be of order $(1.2, 1.1, 1.1, 3.9) \times 10^{-5}$, respectively. They are readily accessible to the B factories.

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I. INTRODUCTION

There are several salient features of the charmless three-body baryonic B decays $B \to \mathbf{B}\bar{\mathbf{B}}'M$ with $\mathbf{B}$ being a baryon. First, a peak near the threshold area of the dibaryon invariant mass spectrum has been observed in many baryonic B decays [1–10]. The so-called threshold effect indicates that the B meson is preferred to decay into a baryon-antibaryon pair with low invariant mass accompanied by a fast recoil meson. Second, none of the two-body charmless baryonic B decays has been observed so far and the present limit on their branching ratios has been pushed to the level of 10^{-7} [11,12]. This means that three-body final states usually have rates larger than their two-body counterparts; that is, $\Gamma(B \to \mathbf{B}\bar{\mathbf{B}}'M) > \Gamma(B \to \mathbf{B}\bar{\mathbf{B}}')$. 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 (for a review, see [13]). The configuration of the two-body decay $B \to \mathbf{B}\bar{\mathbf{B}}'$ is not favorable since its invariant mass is m_B. In $B \to \mathbf{B}\bar{\mathbf{B}}'M$ decays, the effective mass of the baryon pair is reduced as the emitted meson can carry away a large amount of the energies released in baryonic B decays. Third, in the dibaryon rest frame the outgoing meson tends to have a correlation with the baryon or the antibaryon. Experimentally, the correlation can be studied by measuring the Dalitz plot asymmetry or the angular distribution of the baryon or the antibaryon in the dibaryon rest frame. For example, the proton in the dibaryon rest frame of $B^- \to p\bar{p}K^-$ is found to prefer to move collinearly with its meson partner [3,6,8], whereas the proton and π^- in $B^- \to p\bar{p}\pi^-$ move back to back most of the time [8].

While various theoretical ideas [13–20] have been put forward to explain the low mass threshold enhancement, this effect can be understood in terms of a simple short-distance picture [21]. One energetic $q\bar{q}$ pair must be emitted back to back by a hard gluon in order to produce a baryon and an antibaryon in the two-body decay. This hard gluon is highly off mass shell and hence the two-body decay amplitude is suppressed by order of α_s/q^2. In the three-body baryonic B decays, a possible configuration is that the $\mathbf{B}\bar{\mathbf{B}}'$ pair is emitted collinearly against the meson. The quark and antiquark pair emitted from a gluon is moving nearly in the same direction. Since this gluon is close to its mass shell, the corresponding configuration is not subject to the short-distance suppression. This implies that the dibaryon pair tends to have a small invariant mass.

While the above-mentioned short-distance picture predicts a correct angular distribution pattern for the decays $B^- \to p\bar{p}\pi^-$, $B^- \to \Lambda_c^+\bar{p}\pi^+$, and $B^- \to \Lambda\bar{p}\gamma$ [4,8,22], it fails to explain the angular correlation observed in $B^- \to p\bar{p}K^-$ [3,8] and $B^- \to \Lambda\bar{p}\pi^-$ [7]. The intuitive argument that the K^- in the $p\bar{p}$ rest frame is expected to emerge parallel to $\bar{p}$ is not borne out by experiment [13]. Likewise, the naive argument that the pion has no preference for its correlation with the Λ or the $\bar{p}$ in the decay $B^- \to \Lambda\bar{p}\pi^-$ is ruled out by the new Belle experiment [7] in which a strong correlation between the Λ and the pion is seen. Therefore, although the short-distance picture appears to describe the global features of the three-body baryonic B decays, it cannot explain the angular correlation enigma encountered in the penguin-dominated processes $B^- \to p\bar{p}K^-$ and $B^- \to \Lambda\bar{p}\pi^-$.

The study of the charmful baryonic B decays $B \to \mathbf{B}\bar{\mathbf{B}}'M_c$ may help improve our understanding of the underlying mechanism for the threshold enhancement and the angular distribution in three-body decays. First, the aforementioned three features also manifest themselves in $B \to \mathbf{B}\bar{\mathbf{B}}'M_c$. An enhancement at the low dibaryon mass has been seen, for example, in the decay $B \to p\bar{p}D^{(*)}$ [23,24], while the Dalitz plot of $B \to p\bar{p}D^{(*)}$ [24] with asymmetric distributions signals a nonzero angular distribution asymmetry. Therefore, $B \to \mathbf{B}\bar{\mathbf{B}}'M_c$ is a mirror reflecting the phenomena observed in $B \to \mathbf{B}\bar{\mathbf{B}}'M$. Second, since the decay rates of $B \to \mathbf{B}\bar{\mathbf{B}}'M_c$ are 10–1000 times larger than that of $B \to \mathbf{B}\bar{\mathbf{B}}'M$, in principle this will render the detection of angular distributions and Dalitz asymmetries in the latter easier. Theoretically, the relevant topologic quark diagrams for charmful decays are simpler than the charmless ones as the latter receive not only tree but also penguin contributions. In this work, we shall therefore focus on $B \to \mathbf{B}\bar{\mathbf{B}}'M_c$ with $M_c = D^{(*)}$ or J/Ψ, and our first task is to see if we can explain the branching fractions.

This paper is organized as follows. The experimental data for the decay rates of $B \to \mathbf{B}\bar{\mathbf{B}}' M_c$ are collected in Sec. II. The formulism is set in Sec. III with special attention to various transition form factors. We then proceed to have a numerical analysis in Sec. IV and discuss some physical results in Sec. V. Section VI contains our conclusions.

II. EXPERIMENTAL DATA

In this section, we summarize the measured branching ratios for the three-body charmful baryonic B decays $B \to \mathbf{B}\bar{\mathbf{B}}' M_c$:

$$\begin{aligned} \mathrm{Br}(\bar{B}^0 \to n\bar{p}D^{*+}) &= (14.5^{+3.4}_{-3.0} \pm 2.7) \times 10^{-4} \quad (\text{CLEO}), \\ \mathrm{Br}(\bar{B}^0 \to p\bar{p}D^0) &= (1.13 \pm 0.06 \pm 0.08) \times 10^{-4} \quad (\text{BABAR}), \\ &= (1.18 \pm 0.15 \pm 0.16) \times 10^{-4} \quad (\text{Belle}), \\ \mathrm{Br}(\bar{B}^0 \to p\bar{p}D^{*0}) &= (1.01 \pm 0.10 \pm 0.09) \times 10^{-4} \quad (\text{BABAR}), \\ &= (1.20^{+0.33}_{-0.29} \pm 0.21) \times 10^{-4} \quad (\text{Belle}), \\ \mathrm{Br}(\bar{B}^0 \to \Lambda\bar{p}D_s^+) &= (2.9 \pm 0.7 \pm 0.5 \pm 0.4) \times 10^{-5} \quad (\text{Belle}), \\ \mathrm{Br}(B^- \to p\bar{p}D^-) &< 0.15 \times 10^{-4} \quad (90\%\ \text{C.L.}) \quad (\text{Belle}), \\ \mathrm{Br}(B^- \to p\bar{p}D^{*-}) &< 0.15 \times 10^{-4} \quad (90\%\ \text{C.L.}) \quad (\text{Belle}), \end{aligned} \tag{1}$$

for $M_c = D^{(*)}$ [23–26], and

$$\begin{aligned} \mathrm{Br}(B^- \to \Lambda\bar{p}J/\Psi) &= (11.6^{+7.4+4.2}_{-5.3-1.8}) \times 10^{-6} \quad (\text{BABAR}), \\ &= (11.6 \pm 2.8^{+1.8}_{-2.3}) \times 10^{-6} \quad (\text{Belle}), \\ \mathrm{Br}(B^- \to \Sigma^0\bar{p}J/\Psi) &< 1.1 \times 10^{-5} \quad (90\%\ \text{C.L.}) \quad (\text{Belle}), \\ \mathrm{Br}(\bar{B}^0 \to p\bar{p}J/\Psi) &< 1.9 \times 10^{-6} \quad (90\%\ \text{C.L.}) \quad (\text{BABAR}), \\ &< 8.3 \times 10^{-7} \quad (90\%\ \text{C.L.}) \quad (\text{Belle}), \end{aligned} \tag{2}$$

for $M_c = J/\psi$ [27,28]. Among various measured decay modes, $\bar{B}^0 \to n\bar{p}D^{*+}$ was first observed by CLEO in 2001 [25]. Note that the decays $\bar{B}^0 \to p\bar{p}D^0$ and $\bar{B}^0 \to p\bar{p}D^{*0}$ have similar results in rates.

III. FORMALISM

The effective Hamiltonian responsible for charmful baryonic B decays reads [29]

$$\mathcal{H}_{\mathrm{eff}} = \frac{G_F}{\sqrt{2}} V_{cb} V^*_{qq'} (c_1 O_1 + c_2 O_2) + \mathrm{H.c.}, \tag{3}$$

where V_{ij} are the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements and c_i the Wilson coefficients. The four-quark operators O_i are defined by

$$\begin{aligned} O_1 &= (\bar{q}'_\alpha q_\alpha)_{V-A} (\bar{c}_\beta b_\beta)_{V-A}, \\ O_2 &= (\bar{q}'_\alpha q_\beta)_{V-A} (\bar{c}_\beta b_\alpha)_{V-A}, \end{aligned} \tag{4}$$

where $(\bar{q}'_i q_j)_{V-A}$ denotes $\bar{q}'_i \gamma_\mu (1 - \gamma_5) q_j$ with $q' = d$, and s, $q = u$, c, α, and β are the color indices.

To proceed, we shall adopt the generalized factorization approach [30,31] in which the vertex and penguin corrections to hadronic matrix elements of four-quark operators are absorbed in the effective Wilson coefficients c_i^{eff} so that the μ dependence in transition matrix elements is smeared out:

$$\sum c_i(\mu) \langle Q_i(\mu) \rangle \to \sum c_i^{\mathrm{eff}} \langle Q_i \rangle_{\mathrm{VIA}}, \tag{5}$$

where the subscript VIA means that the hadronic matrix element is evaluated under the vacuum insertion approximation. This approach has been well applied to the study of three-body baryonic B decays [13,17–19,32–44].

Under the factorization approximation, the decay amplitudes can be classified into three different categories: the current-type (class-I), the transition-type (class-II), and the hybrid-type (class-III) amplitudes. The class-I current-type amplitudes which proceed via a color-allowed, external

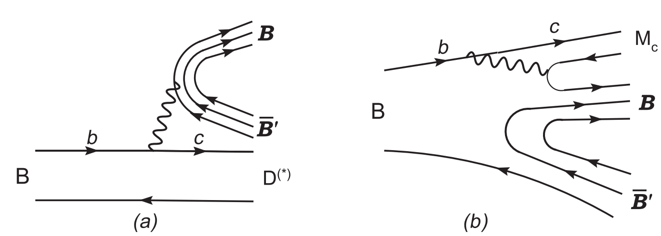

FIG. 1. Two types of the $B \to \mathbf{B}\bar{\mathbf{B}}' M_c$ decay process: (a) current type and (b) transition type.

W-emission diagram as depicted in Fig. 1(a) are given by

$$\mathcal{A}_{\mathcal{C}}(B \to \mathbf{B}\bar{\mathbf{B}}' D^{(*)}) = \frac{G_F}{\sqrt{2}} V_{cb} V^*_{uq'} a_1^{D^{(*)}} \langle \mathbf{B}\bar{\mathbf{B}}' | (\bar{q}' u)_{V-A} | 0 \rangle \times \langle D^{(*)} | (\bar{c} b)_{V-A} | B \rangle, \quad (6)$$

with $a_1^{D^{(*)}}$ to be specified later. By fixing $q = u$ to avoid the charm baryon production, all possible decay modes in this category are

$$\begin{aligned} &\bar{B}^0 \to \{n\bar{p}, \Sigma^-\bar{\Lambda}, \Lambda\bar{\Sigma}^+, \Sigma^-\bar{\Sigma}^0, \Sigma^0\bar{\Sigma}^+, \Xi^-\bar{\Xi}^0\} D^{(*)+} \quad \text{for } q' = d, \\ &\bar{B}^0 \to \{\Lambda\bar{p}, \Sigma^0\bar{p}, \Sigma^-\bar{n}, \Xi^-\bar{\Sigma}^0, \Xi^-\bar{\Lambda}, \Xi^0\bar{\Sigma}^+\} D^{(*)+} \quad \text{for } q' = s. \end{aligned} \quad (7)$$

The class-II transition-type amplitude via the color-suppressed internal W emission diagram [Fig. 1(b)] reads

$$\mathcal{A}_{\mathcal{T}}(B \to \mathbf{B}\bar{\mathbf{B}}' M_c) = \frac{G_F}{\sqrt{2}} V_{cb} V^*_{qq'} a_2^{M_c} \langle M_c | (\bar{c} q)_{V-A} | 0 \rangle \langle \mathbf{B}\bar{\mathbf{B}}' | (\bar{q}' b)_{V-A} | B \rangle, \quad (8)$$

with $a_2^{M_c}$ to be given later. The possible decay modes in this class are

$$\begin{aligned} &B^- \to \{n\bar{p}, \Sigma^-\bar{\Lambda}, \Lambda\bar{\Sigma}^+, \Sigma^-\bar{\Sigma}^0, \Sigma^0\bar{\Sigma}^+, \Xi^-\bar{\Xi}^0\} J/\Psi \quad \text{for } q' = d, \\ &B^- \to \{\Lambda\bar{p}, \Sigma^0\bar{p}, \Sigma^-\bar{n}, \Xi^-\bar{\Sigma}^0, \Xi^-\bar{\Lambda}, \Xi^0\bar{\Sigma}^+\} J/\Psi \quad \text{for } q' = s, \end{aligned} \quad (9)$$

for the charged B system, and

$$\begin{aligned} &\bar{B}^0 \to \{p\bar{p}, \Lambda\bar{\Lambda}, \Sigma^+\bar{\Sigma}^+, \Xi^-\bar{\Xi}^-, \Lambda\bar{\Sigma}^0, n\bar{n}, \Sigma^0\bar{\Sigma}^0, \Sigma^-\bar{\Sigma}^-, \Xi^0\bar{\Xi}^0, \Sigma^0\bar{\Lambda}\} D^{(*)0}(J/\Psi) \quad \text{for } q' = d, \\ &\bar{B}^0 \to \{\Lambda\bar{n}, \Sigma^0\bar{n}, \Sigma^+\bar{p}, \Xi^0\bar{\Sigma}^0, \Xi^0\bar{\Lambda}, \Xi^-\bar{\Xi}^-\} D^{(*)0}(J/\Psi) \quad \text{for } q' = s, \end{aligned} \quad (10)$$

for the neutral B system. The class-III hybrid-type amplitudes which consist of both the current-type and the transition-type transitions are given by

$$\mathcal{A}_{\mathcal{H}}(B \to \mathbf{B}\bar{\mathbf{B}}' D^{(*)}) = \frac{G_F}{\sqrt{2}} V_{cb} V^*_{uq'} \{ a_1^{D^{(*)}} \langle \mathbf{B}\bar{\mathbf{B}}' | (\bar{q}' u)_{V-A} | 0 \rangle \langle D^{(*)} | (\bar{c} b)_{V-A} | B \rangle + a_2^{D^{(*)}} \langle D^{(*)} | (\bar{c} u)_{V-A} | 0 \rangle \langle \mathbf{B}\bar{\mathbf{B}}' | (\bar{q}' b)_{V-A} | B \rangle \}. \quad (11)$$

The allowed modes are

$$\begin{aligned} &B^- \to \{n\bar{p}, \Sigma^-\bar{\Lambda}, \Lambda\bar{\Sigma}^+, \Sigma^-\bar{\Sigma}^0, \Sigma^0\bar{\Sigma}^+, \Xi^-\bar{\Xi}^0\} D^{(*)0} \quad \text{for } q' = d, \\ &B^- \to \{\Lambda\bar{p}, \Sigma^0\bar{p}, \Sigma^-\bar{n}, \Xi^-\bar{\Sigma}^0, \Xi^-\bar{\Lambda}, \Xi^0\bar{\Sigma}^+\} D^{(*)0} \quad \text{for } q' = s. \end{aligned} \quad (12)$$

We now can have a qualitative understanding of the data in (1) and (2). The nonobservation of $B^- \to p\bar{p}D^{(*)-}$ is due to the fact that it proceeds via $b \to u\bar{c}d$ at the quark level. This leads to a suppression of $|V_{ub}V^*_{cd}/V_{cb}V^*_{ud}|^2 \simeq 10^{-4}$ compared to its neutral partner. Likewise, it is expected that $\mathrm{Br}(\bar{B}^0 \to p\bar{p}J/\Psi) = |V_{cd}/V_{cs}|^2 \mathrm{Br}(B^- \to \Lambda\bar{p}J/\Psi) \simeq 10^{-7}$, consistent with the experimental upper bound for this decay mode. As for the comparable decay rates for $\bar{B}^0 \to p\bar{p}D^{*0}$ and $\bar{B}^0 \to p\bar{p}D^0$, it has to do with the baryonic form factors, which we are going to elaborate on later.

A. Decay constants and form factors

To proceed, we need the information of the decay constants and form factors relevant for charmful baryonic B decays. Decay constants are defined by

$$\begin{aligned} \langle P_c | \bar{q}_1 \gamma_\mu \gamma_5 q_2 | 0 \rangle &= -i f_{P_c} p_\mu, \\ \langle V_c | \bar{q}_1 \gamma_\mu q_2 | 0 \rangle &= m_{V_c} f_{V_c} \varepsilon^*_\mu, \end{aligned} \quad (13)$$

for the pseudoscalar P_c and the vector meson V_c. The B to $D^{(*)}$ transition form factors are parametrized as [45]

$$\langle D|\bar{c}\gamma^\mu b|B\rangle = \left[(p_B + p_D)^\mu - \frac{m_B^2 - m_D^2}{t}q^\mu\right]F_1^{BD}(t) + \frac{m_B^2 - m_D^2}{t}q^\mu F_0^{BD}(t),$$
$$\langle D^*|\bar{c}\gamma_\mu b|B\rangle = \epsilon_{\mu\nu\alpha\beta}\varepsilon^{*\nu}p_B^\alpha p_{D^*}^\beta \frac{2V_1^{BD^*}(t)}{m_B + m_{D^*}},$$
$$\langle D^*|\bar{c}\gamma_\mu\gamma_5 b|B\rangle = i\left[\varepsilon_\mu^* - \frac{\varepsilon^*\cdot q}{t}q_\mu\right](m_B + m_{D^*})A_1^{BD^*}(t) + i\frac{\varepsilon^*\cdot q}{t}q_\mu(2m_{D^*})A_0^{BD^*}(t) - i\left[(p_B + p_{D^*})_\mu - \frac{m_B^2 - m_{D^*}^2}{t}q_\mu\right](\varepsilon^*\cdot q)\frac{A_2^{BD^*}(t)}{m_B + m_{D^*}}, \tag{14}$$

where $t \equiv q^2$ with $q = p_B - p_{D^*} = p_{\mathbf{B}} + p_{\bar{\mathbf{B}}'}$.

For the dibaryon creation, we write

$$\langle \mathbf{B}\bar{\mathbf{B}}'|\bar{q}_1\gamma_\mu q_2|0\rangle = \bar{u}(p_{\mathbf{B}})\left\{F_1(t)\gamma_\mu + \frac{F_2(t)}{m_{\mathbf{B}} + m_{\bar{\mathbf{B}}'}}i\sigma_{\mu\nu}q_\mu\right\}v(p_{\bar{\mathbf{B}}'}),$$
$$= \bar{u}(p_{\mathbf{B}})\left\{[F_1(t) + F_2(t)]\gamma_\mu + \frac{F_2(t)}{m_{\mathbf{B}} + m_{\bar{\mathbf{B}}'}}(p_{\bar{\mathbf{B}}'} - p_{\mathbf{B}})_\mu\right\}v(p_{\bar{\mathbf{B}}'}), \tag{15}$$
$$\langle \mathbf{B}\bar{\mathbf{B}}'|\bar{q}_1\gamma_\mu\gamma_5 q_2|0\rangle = \bar{u}(p_{\mathbf{B}})\left\{g_A(t)\gamma_\mu + \frac{h_A(t)}{m_{\mathbf{B}} + m_{\bar{\mathbf{B}}'}}q_\mu\right\}\gamma_5 v(p_{\bar{\mathbf{B}}'}),$$

where $u(v)$ is the (anti-)baryon spinor, and $F_{1,2}$, g_A, h_A are timelike baryonic form factors. Note that there are two additional form factors in the form of $\bar{u}q_\mu v$ and $\bar{u}\sigma_{\mu\nu}q_\nu\gamma_5 v$. However, since we assume SU(3) flavor symmetry, we can neglect these two form factors as they vanish for conserved currents. The asymptotic behavior of form factors is governed by the pQCD counting rules [46–48]. In the large t limit, the momentum dependence of the form factors $F_1(t)$ and $g_A(t)$ behaves as $1/t^2$ as there are two hard gluon exchanges between the valence quarks. More precisely, in the $t \to \infty$ limit

$$F_1(t) = \frac{C_{F_1}}{t^2}\left[\ln\left(\frac{t}{\Lambda_0^2}\right)\right]^{-\gamma}, \qquad g_A(t) = \frac{C_{g_A}}{t^2}\left[\ln\left(\frac{t}{\Lambda_0^2}\right)\right]^{-\gamma}, \tag{16}$$

where $\gamma = 2 + 4/(3\beta) = 2.148$ with β being the QCD β function and $\Lambda_0 = 0.3$ GeV. In the asymptotic $t \to \infty$ limit, both $F_2(t)$ and $h_A(t)$ have an extra $1/t$ dependence relative to F_1 and g_A owing to a mass insertion at the quark line [38,49,50]. However, the form factor h_A is related to g_A by the relation

$$h_A = -\frac{(m_{\mathbf{B}} + m_{\bar{\mathbf{B}}'})^2}{t}g_A, \tag{17}$$

through the equation of motion. Hence, in ensuing numerical analysis we will keep $h_A(t)$ and neglect $F_2(t)$. Under the SU(3) flavor and SU(2) spin symmetries [48], the parameters C_{F_1} and C_{g_A} appearing in $0 \to \mathbf{B}\bar{\mathbf{B}}'$ transitions are no longer independent but are related to each other through the two reduced parameters $C_{\|}$ and $C_{\bar{\|}}$ (see Table I).

For the three-body transition $B \to \mathbf{B}\bar{\mathbf{B}}'$, its most general expression reads

$$\langle \mathbf{B}\bar{\mathbf{B}}'|\bar{q}'\gamma_\mu b|B\rangle = i\bar{u}(p_{\mathbf{B}})[g_1\gamma_\mu + g_2 i\sigma_{\mu\nu}p^\nu + g_3 p_\mu + g_4(p_{\bar{\mathbf{B}}'} + p_{\mathbf{B}})_\mu + g_5(p_{\bar{\mathbf{B}}'} - p_{\mathbf{B}})_\mu] \times \gamma_5 v(p_{\bar{\mathbf{B}}'}),$$
$$\langle \mathbf{B}\bar{\mathbf{B}}'|\bar{q}'\gamma_\mu\gamma_5 b|B\rangle = i\bar{u}(p_{\mathbf{B}})[f_1\gamma_\mu + f_2 i\sigma_{\mu\nu}p^\nu + f_3 p_\mu + f_4(p_{\bar{B}'} + p_{\mathbf{B}})_\mu + f_5(p_{\bar{\mathbf{B}}'} - p_{\mathbf{B}})_\mu] \times v(p_{\bar{\mathbf{B}}'}), \tag{18}$$

with $p = p_B - p_{\mathbf{B}} - p_{\bar{\mathbf{B}}'}$. In principle, the form factors f_i and g_i depend on not only the invariant mass t of the dibaryon, but also the invariant mass (i.e., the variable u

TABLE I. The parameters C_{F_1} and C_{g_A} in Eq. (16), expressed in terms of the combination of $C_{\|}$ and $C_{\bar{\|}}$, where the upper (lower) sign is for C_{F_1} (C_{g_A}).

$0 \to$	$C_{F_1}(C_{g_A})$
$n\bar{p}$	$\frac{4}{3}C_{\|} \mp \frac{1}{3}C_{\bar{\|}}$
$\Sigma^-\bar{\Lambda}$	$\frac{1}{\sqrt{6}}C_{\|} \mp \frac{1}{\sqrt{6}}C_{\bar{\|}}$
$\Lambda\bar{\Sigma}^+$	$\frac{1}{\sqrt{6}}C_{\|} \mp \frac{1}{\sqrt{6}}C_{\bar{\|}}$
$\Sigma^-\bar{\Sigma}^0$	$\frac{5}{3\sqrt{2}}C_{\|} \pm \frac{1}{3\sqrt{2}}C_{\bar{\|}}$
$\Sigma^0\bar{\Sigma}^+$	$\frac{-5}{3\sqrt{2}}C_{\|} \mp \frac{1}{3\sqrt{2}}C_{\bar{\|}}$
$\Xi^-\bar{\Xi}^0$	$\frac{-1}{3}C_{\|} \mp \frac{2}{3}C_{\bar{\|}}$
$\Lambda\bar{p}$	$-\sqrt{\frac{3}{2}}C_{\|}$
$\Sigma^0\bar{p}$	$\frac{-1}{3\sqrt{2}}C_{\|} \mp \frac{2}{3\sqrt{2}}C_{\bar{\|}}$
$\Sigma^-\bar{n}$	$\frac{-1}{3}C_{\|} \mp \frac{2}{3}C_{\bar{\|}}$
$\Xi^-\bar{\Sigma}^0$	$\frac{4}{3\sqrt{2}}C_{\|} \mp \frac{1}{3\sqrt{2}}C_{\bar{\|}}$
$\Xi^-\bar{\Lambda}$	$\frac{2}{\sqrt{6}}C_{\|} \pm \frac{1}{\sqrt{6}}C_{\bar{\|}}$
$\Xi^0\bar{\Sigma}^+$	$\frac{4}{3}C_{\|} \mp \frac{1}{3}C_{\bar{\|}}$

or s) of one of the baryons and the emitted meson. Indeed, such a momentum dependence has been studied in the framework of the pole model in which some intermediate pole contributions to the three-body transition of $B \to \mathbf{B}\bar{\mathbf{B}}'$ are considered [13,17–19,32]. However, since the momentum dependence of the transition form factors on the variable s or u is poorly known, one approach which is often employed in the literature is that one first parametrizes them in the form of a power series of the dibaryon invariant mass t and assumes that the dependence on other variables may be lumped into the coefficients of $1/t$ expansion [36,38–41,51]. This is the approach we will adapt in this work. Since three gluons are needed to induce the $B \to \mathbf{B}\bar{\mathbf{B}}'$ transition, two for producing the baryon pair and one for kicking the spectator quark in the B meson, pQCD counting rules imply that to the leading order

TABLE II. The parameters D_{g_1}, D_{f_1}, D_{g_i}, and D_{f_i} in $B \to \mathbf{B}\bar{\mathbf{B}}'$ transition expressed in terms of the reduced parameters $D_{\parallel}$, $D_{\bar{\parallel}}$, $D^i_{\parallel}$ with $i = 2, 3, 4, 5$, where the upper (lower) sign is for D_{g_1} (D_{f_1}).

$B^- \to$	$D_{g_1}(D_{f_1})$	D_{g_i}, $-D_{f_i}$
$n\bar{p}$	$\frac{4}{3}D_{\parallel} \pm \frac{1}{3}D_{\bar{\parallel}}$	$\frac{5}{3}D^i_{\parallel}$
$\Sigma^-\bar{\Lambda}$	$\frac{1}{\sqrt{6}}D_{\parallel} \pm \frac{1}{\sqrt{6}}D_{\bar{\parallel}}$	$\sqrt{\frac{2}{3}}D^i_{\parallel}$
$\Lambda\bar{\Sigma}^+$	$\frac{1}{\sqrt{6}}D_{\parallel} \pm \frac{1}{\sqrt{6}}D_{\bar{\parallel}}$	$\sqrt{\frac{2}{3}}D^i_{\parallel}$
$\Sigma^-\bar{\Sigma}^0$	$\frac{5}{3\sqrt{2}}D_{\parallel} \mp \frac{1}{3\sqrt{2}}D_{\bar{\parallel}}$	$\frac{2\sqrt{2}}{3}D^i_{\parallel}$
$\Sigma^0\bar{\Sigma}^+$	$-\frac{5}{3\sqrt{2}}D_{\parallel} \pm \frac{1}{3\sqrt{2}}D_{\bar{\parallel}}$	$-\frac{2\sqrt{2}}{3}D^i_{\parallel}$
$\Xi^-\bar{\Xi}^0$	$-\frac{1}{3}D_{\parallel} \pm \frac{2}{3}D_{\bar{\parallel}}$	$\frac{1}{3}D^i_{\parallel}$
$\Lambda\bar{p}$	$-\sqrt{\frac{3}{2}}D_{\parallel}$	$-\sqrt{\frac{3}{2}}D^i_{\parallel}$
$\Sigma^0\bar{p}$	$-\frac{1}{3\sqrt{2}}D_{\parallel} \pm \frac{2}{3\sqrt{2}}D_{\bar{\parallel}}$	$\frac{1}{3\sqrt{2}}D^i_{\parallel}$
$\Sigma^-\bar{n}$	$-\frac{1}{3}D_{\parallel} \pm \frac{2}{3}D_{\bar{\parallel}}$	$\frac{1}{3}D^i_{\parallel}$
$\Xi^-\bar{\Sigma}^0$	$\frac{4}{3\sqrt{2}}D_{\parallel} \pm \frac{1}{3\sqrt{2}}D_{\bar{\parallel}}$	$\frac{5}{3\sqrt{2}}D^i_{\parallel}$
$\Xi^-\bar{\Lambda}$	$\frac{2}{\sqrt{6}}D_{\parallel} \mp \frac{1}{\sqrt{6}}D_{\bar{\parallel}}$	$\frac{1}{\sqrt{6}}D^i_{\parallel}$
$\Xi^0\bar{\Sigma}^+$	$\frac{4}{3}D_{\parallel} \pm \frac{1}{3}D_{\bar{\parallel}}$	$\frac{5}{3}D^i_{\parallel}$

$\bar{B}^0 \to$	$D_{g_1}(D_{f_1})$	D_{g_i}, $-D_{f_i}$
$p\bar{p}$, $\Xi^-\bar{\Xi}^-$	$\frac{1}{3}D_{\parallel} \mp \frac{2}{3}D_{\bar{\parallel}}$	$-\frac{1}{3}D^i_{\parallel}$
$n\bar{n}$, $\Sigma^-\bar{\Sigma}^-$	$\frac{5}{3}D_{\parallel} \mp \frac{1}{3}D_{\bar{\parallel}}$	$\frac{4}{3}D^i_{\parallel}$
$\Lambda\bar{\Lambda}$	$\frac{1}{2}D_{\parallel} \mp \frac{1}{2}D_{\bar{\parallel}}$	0
$\Sigma^0\bar{\Sigma}^0$	$\frac{5}{6}D_{\parallel} \mp \frac{1}{6}D_{\bar{\parallel}}$	$\frac{2}{3}D^i_{\parallel}$
$\Sigma^+\bar{\Sigma}^+$, $\Xi^0\bar{\Xi}^0$	0	0
$\Lambda\bar{\Sigma}^0$, $\Sigma^0\bar{\Lambda}$	$-\frac{1}{2\sqrt{3}}D_{\parallel} \mp \frac{1}{2\sqrt{3}}D_{\bar{\parallel}}$	$-\frac{1}{\sqrt{3}}D^i_{\parallel}$
$\Lambda\bar{n}$	$-\sqrt{\frac{3}{2}}D_{\parallel}$	$-\sqrt{\frac{3}{2}}D^i_{\parallel}$
$\Sigma^0\bar{n}$	$\frac{1}{3\sqrt{2}}D_{\parallel} \mp \frac{2}{3\sqrt{2}}D_{\bar{\parallel}}$	$-\frac{1}{3\sqrt{2}}D^i_{\parallel}$
$\Sigma^+\bar{p}$	$-\frac{1}{3}D_{\parallel} \pm \frac{2}{3}D_{\bar{\parallel}}$	$\frac{1}{3}D^i_{\parallel}$
$\Xi^0\bar{\sigma}^0$	$-\frac{4}{3\sqrt{2}}D_{\parallel} \mp \frac{1}{3\sqrt{2}}D_{\bar{\parallel}}$	$-\frac{5}{3\sqrt{2}}D^i_{\parallel}$
$\Xi^0\bar{\Lambda}$	$\frac{2}{\sqrt{6}}D_{\parallel} \mp \frac{1}{\sqrt{6}}D_{\bar{\parallel}}$	$\frac{1}{\sqrt{6}}D^i_{\parallel}$
$\Xi^-\bar{\Sigma}^-$	$\frac{4}{3}D_{\parallel} \pm \frac{1}{3}D_{\bar{\parallel}}$	$\frac{5}{3}D^i_{\parallel}$

$$f_i(t) = \frac{D_{f_i}}{t^3}, \qquad g_i(t) = \frac{D_{g_i}}{t^3}. \tag{19}$$

Just as the previous case for vacuum to the dibaryon transition, under the SU(3) flavor and SU(2) spin symmetries, the parameters D_{g_1}, D_{f_1}, D_{g_i}, and D_{f_i} can be expressed in terms of the reduced parameters $D_{\parallel}$, $D_{\bar{\parallel}}$, and $D^i_{\parallel}$ (see Table II), and the derivation is presented in the Appendix. Interestingly, the decays $\bar{B}^0 \to \Sigma^+\bar{\Sigma}^+$ and $\Xi^0\bar{\Xi}^0$ are prohibited by SU(2) spin symmetry.

Since many of the three-body baryonic B decays show an enhancement at the low dibaryon invariant mass squared t, low values of t produce the dominant contributions to these decays. Recently, *BABAR* has measured the phase-space corrected dibaryon invariant mass distributions for various $B \to D^{(*)}p\bar{p}$ decays [24]. This amounts to measuring the squared amplitude $|A|^2$ versus t. It turns out that the t dependence of $|A|^2$ can be represented sufficiently by a single $1/t^n$ behavior. Therefore, it is justified to apply the pQCD counting rules valid at large t to form factors at small values of t.

IV. NUMERICAL ANALYSIS

We need to specify various input parameters for a numerical analysis. For the CKM matrix elements, we use the Wolfenstein parameters $A = 0.807 \pm 0.018$, $\lambda = 0.2265 \pm 0.0008$, $\bar{\rho} = 0.141$, and $\bar{\eta} = 0.343$ [52]. For the decay constants, we use [53,54]

$$(f_D, f_{D^*}, f_{J/\Psi}) = (0.22, 0.23, 0.41)\ \text{GeV}. \tag{20}$$

Following [55], the momentum dependence of $B \to D^{(*)}$ transition form factors in Eq. (14) is parametrized as

$$F(q^2) = \frac{F(0)}{1 - a(q^2/m_B^2) + b(q^2/m_B^2)^2}. \tag{21}$$

We shall use the parameters a, b, and $F(0)$ obtained in [55] (see Table III).

For the parameters $C_{\parallel}$ and $C_{\bar{\parallel}}$ in Eqs. (15) and (16) and Table I, we use the data of $e^+e^- \to p\bar{p}$, $n\bar{n}$ [56,57] to determine their magnitudes and the decay rate of $\bar{B}^0 \to n\bar{p}D^{*+}$ to fix their relative sign

$$(C_{\parallel}, C_{\bar{\parallel}}) = (67.9 \pm 1.4, -216.9 \pm 23.5)\ \text{GeV}^4, \tag{22}$$

with $\chi^2/\text{d.o.f} = 1.6$. As for the parameters $D_{\parallel}$ and $D_{\bar{\parallel}}$ in Eqs. (18) and (19) and Table II, we employ the observed rates of $\bar{B}^0 \to p\bar{p}D^0$, $B^- \to p\bar{p}K^{*-}$, $\bar{B}^0 \to p\bar{p}K^{*0}$, and

TABLE III. Various $B \to D^{(*)}$ form factors taken from [55].

$B \to D^{(*)}$	F_1	F_0	V_1	A_0	A_1	A_2
$F(0)$	0.67	0.67	0.75	0.64	0.63	0.61
a	1.25	0.65	1.29	1.30	0.65	1.14
b	0.39	0.00	0.45	0.31	0.02	0.52

TABLE IV. Branching ratios for the current-type $\bar{B}^0 \to \mathbf{B}\bar{\mathbf{B}}' D^{(*)+}$ decays. The results for D^* production are shown in parentheses.

$\mathbf{B}\bar{\mathbf{B}}'$	$\mathrm{Br}(\bar{B}^0 \to \mathbf{B}\bar{\mathbf{B}}' D^{(*)+}) \times 10^4$
$n\bar{p}$	5.2 ± 0.5 (14.5 ± 1.4)
$\Sigma^-\bar{\Lambda}$	0.3 ± 0.1 (1.3 ± 0.2)
$\Lambda\bar{\Sigma}^+$	0.3 ± 0.1 (1.3 ± 0.2)
$\Sigma^-\bar{\Sigma}^0$	0.05 ± 0.01 (0.23 ± 0.02)
$\Sigma^0\bar{\Sigma}^+$	0.05 ± 0.01 (0.23 ± 0.02)
$\Xi^-\bar{\Xi}^0$	0.07 ± 0.02 (0.4 ± 0.1)

$\mathbf{B}\bar{\mathbf{B}}'$	$\mathrm{Br}(\bar{B}^0 \to \mathbf{B}\bar{\mathbf{B}}' D^{(*)+}) \times 10^6$
$\Lambda\bar{p}$	3.4 ± 0.2 (11.9 ± 0.5)
$\Sigma^0\bar{p}$	2.8 ± 0.7 (10.5 ± 2.7)
$\Sigma^-\bar{n}$	5.5 ± 1.4 (20.9 ± 5.4)
$\Xi^-\bar{\Sigma}^0$	0.6 ± 0.1 (3.0 ± 0.3)
$\Xi^-\bar{\Lambda}$	0.2 ± 0.1 (1.3 ± 0.4)
$\Xi^0\bar{\Sigma}^+$	1.2 ± 0.1 (6.2 ± 0.6)

$B^- \to p\bar{p}\pi^-$ [2,3,8–10,23,24] in conjunction with the measured angular distribution of the last decay mode to obtain

$$\begin{aligned}(D_\parallel, D_{\bar\parallel}) &= (67.7 \pm 16.3, -280.0 \pm 35.9)\ \mathrm{GeV}^5,\\ (D_\parallel^2, D_\parallel^3, D_\parallel^4, D_\parallel^5) &= (-187.3 \pm 26.6, -840.1 \pm 132.1,\\ &\quad -10.1 \pm 10.8, -157.0 \pm 27.1)\ \mathrm{GeV}^4,\end{aligned} \tag{23}$$

with $\chi^2/\mathrm{d.o.f} = 0.8$. To calculate the decay rates, we use the relation

TABLE V. Same as Table IV except for the transition-type $\bar{B}^0 \to \mathbf{B}\bar{\mathbf{B}}' D^{(*)0}$ decays.

$\mathbf{B}\bar{\mathbf{B}}'$	$\mathrm{Br}(\bar{B}^0 \to \mathbf{B}\bar{\mathbf{B}}' D^{(*)0}) \times 10^4$
$p\bar{p}$	1.1 ± 0.2 (1.0 ± 0.4)
$n\bar{n}$	3.9 ± 1.3 (5.8 ± 2.3)
$\Lambda\bar{\Lambda}$	0.02 ± 0.01 (0.03 ± 0.02)
$\Sigma^0\bar{\Sigma}^0$	0.10 ± 0.03 (0.07 ± 0.03)
$\Sigma^-\bar{\Sigma}^-$	0.4 ± 0.1 (0.3 ± 0.1)
$\Xi^-\bar{\Xi}^-$	0.022 ± 0.004 (0.012 ± 0.005)
$\{^{\Lambda\bar{\Sigma}^0}_{\Sigma^0\bar{\Lambda}}$	0.18 ± 0.05 (0.12 ± 0.04)

$\mathbf{B}\bar{\mathbf{B}}'$	$\mathrm{Br}(\bar{B}^0 \to \mathbf{B}\bar{\mathbf{B}}' D^{(*)0}) \times 10^6$
$\Lambda\bar{n}$	8.2 ± 2.7 (9.8 ± 3.7)
$\Sigma^0\bar{n}$	0.7 ± 0.1 (0.6 ± 0.2)
$\Sigma^+\bar{p}$	1.4 ± 0.3 (1.2 ± 0.4)
$\Xi^0\bar{\Sigma}^0$	1.2 ± 0.4 (0.7 ± 0.3)
$\Xi^0\bar{\Lambda}$	0.09 ± 0.04 (0.10 ± 0.04)
$\Xi^-\bar{\Sigma}^-$	2.2 ± 0.7 (1.4 ± 0.5)

TABLE VI. Same as Table IV except for the hybrid-type $B^- \to \mathbf{B}\bar{\mathbf{B}}' D^{(*)0}$ decays.

$\mathbf{B}\bar{\mathbf{B}}'$	$\mathrm{Br}(B^- \to \mathbf{B}\bar{\mathbf{B}}' D^{(*)0}) \times 10^4$
$n\bar{p}$	13.0 ± 2.9 (46.6 ± 5.3)
$\Sigma^-\bar{\Lambda}$	0.6 ± 0.1 (2.6 ± 0.3)
$\Lambda\bar{\Sigma}^+$	0.6 ± 0.1 (2.6 ± 0.3)
$\Sigma^-\bar{\Sigma}^0$	0.3 ± 0.1 (0.6 ± 0.1)
$\Sigma^0\bar{\Sigma}^+$	0.3 ± 0.1 (0.6 ± 0.1)
$\Xi^-\bar{\Xi}^0$	0.09 ± 0.01 (0.6 ± 0.1)

$\mathbf{B}\bar{\mathbf{B}}'$	$\mathrm{Br}(B^- \to \mathbf{B}\bar{\mathbf{B}}' D^{(*)0}) \times 10^6$
$\Lambda\bar{p}$	11.4 ± 2.6 (32.3 ± 3.2)
$\Sigma^0\bar{p}$	3.3 ± 0.6 (15.2 ± 2.4)
$\Sigma^-\bar{n}$	7.0 ± 1.2 (30.2 ± 4.7)
$\Xi^-\bar{\Sigma}^0$	1.9 ± 0.3 (6.2 ± 0.5)
$\Xi^-\bar{\Lambda}$	0.5 ± 0.1 (1.2 ± 0.2)
$\Xi^0\bar{\Sigma}^+$	3.7 ± 0.7 (12.7 ± 1.1)

$$d\Gamma = \frac{1}{(2\pi)^3}\frac{|\bar{A}|^2}{32M_B^3} dm^2_{\mathbf{B}\bar{\mathbf{B}}'} dm^2_{\bar{\mathbf{B}}'M_c} \tag{24}$$

with $m^2_{\mathbf{B}\bar{\mathbf{B}}'} = (P_{\mathbf{B}} + P_{\bar{\mathbf{B}}'})^2$, $m^2_{\bar{\mathbf{B}}'M_c} = (P_{\bar{\mathbf{B}}'} + P_{M_c})^2$, and the baryon spins have been summed over in $|\bar{A}|^2$. The numerical results for the branching ratios are summarized in Tables IV, V, VI, and VII, where the theoretical errors come from the uncertainties in form factors.

V. DISCUSSION

In the generalized factorization approach, the coefficients a_1 and a_2 appearing in the decay amplitudes $\mathcal{A}_{\mathcal{C}}$, $\mathcal{A}_{\mathcal{T}}$, and $\mathcal{A}_{\mathcal{H}}$ are given by [29–31]

$$a_1 = c_1^{\mathrm{eff}} + \frac{c_2^{\mathrm{eff}}}{3} + c_2^{\mathrm{eff}}\chi_1, \qquad a_2 = c_2^{\mathrm{eff}} + \frac{c_1^{\mathrm{eff}}}{3} + c_1^{\mathrm{eff}}\chi_2, \tag{25}$$

TABLE VII. Branching ratios for the transition-type $B^- \to \mathbf{B}\bar{\mathbf{B}}' J/\Psi$ decays.

$\mathbf{B}\bar{\mathbf{B}}'$	$\mathrm{Br}(B^- \to \mathbf{B}\bar{\mathbf{B}}' J/\Psi) \times 10^6$
$n\bar{p}$	14.4 ± 3.2
$\Sigma^-\bar{\Lambda}$	0.04 ± 0.02
$\Lambda\bar{\Sigma}^+$	0.04 ± 0.02
$\Sigma^-\bar{\Sigma}^0$	0.29 ± 0.09
$\Sigma^0\bar{\Sigma}^+$	0.29 ± 0.09
$\Xi^-\bar{\Xi}^0$	0.37 ± 0.08
$\Lambda\bar{p}$	11.6 ± 2.9
$\Sigma^0\bar{p}$	0.13 ± 0.02
$\Sigma^-\bar{n}$	0.24 ± 0.05
$\Xi^-\bar{\Sigma}^0$	22.9 ± 5.8
$\Xi^-\bar{\Lambda}$	0.52 ± 0.33
$\Xi^0\bar{\Sigma}^+$	44.8 ± 10.8

TABLE VIII. Branching ratios for the transition-type $\bar{B}^0 \to \mathbf{B}\bar{\mathbf{B}}' J/\Psi$ decays.

$\mathbf{B}\bar{\mathbf{B}}'$	$\mathrm{Br}(\bar{B}^0 \to \mathbf{B}\bar{\mathbf{B}}' J/\Psi) \times 10^6$
$p\bar{p}$	1.2 ± 0.2
$n\bar{n}$	6.7 ± 1.7
$\Lambda\bar{\Lambda}$	0.0010 ± 0.0004
$\Sigma^0\bar{\Sigma}^0$	0.13 ± 0.04
$\Sigma^-\bar{\Sigma}^-$	0.54 ± 0.16
$\Xi^-\bar{\Xi}^-$	0.35 ± 0.08
$\{^{\Lambda\bar{\Sigma}^0}_{\Sigma^0\bar{\Lambda}}$	0.09 ± 0.02
$\Lambda\bar{n}$	10.4 ± 2.6
$\Sigma^0\bar{n}$	0.11 ± 0.02
$\Sigma^+\bar{p}$	0.24 ± 0.05
$\Xi^0\bar{\Sigma}^0$	20.6 ± 4.8
$\Xi^0\bar{\Lambda}$	0.5 ± 0.3
$\Xi^-\bar{\Sigma}^-$	42.1 ± 10.4

with $(c_1^{\mathrm{eff}}, c_2^{\mathrm{eff}}) = (1.169, -0.367)$. Since the nonfactorizable effects characterized by the $c_i^{\mathrm{eff}}\chi_j$ terms are not calculable, we will fit them to the data. We see that a_2 is dominated by the nonfactorizable term due to the large cancellation between c_2^{eff} and $c_1^{\mathrm{eff}}/3$. Hence, nonfactorizable contributions to color-suppressed modes are sizable. Factorization works if the parameters a_1 and a_2 are universal; namely, they are channel by channel independent. Empirically, this is supported by the experimental measurements [58]. Two-body mesonic B decays suggest that $a_1 \sim \mathcal{O}(1)$ and $a_2 \sim \mathcal{O}(0.2\text{–}0.3)$ [58,59]. It will become plausible if $a_i^{M_c}$ fitted from the baryonic B decays lie in the aforementioned ranges. Indeed, the results given by

$$a_1^{D^*} = 1.23 \pm 0.19, \qquad a_2^{D^*} = 0.33 \pm 0.04, \qquad a_2^{J/\Psi} = 0.17 \pm 0.03, \tag{26}$$

agree with the above expectation. Hence, we shall assume the validity of factorization in charmful baryonic B decays.

There are several salient features we learn from Tables IV, V, and VI: (i) $\mathcal{B}(\bar{B}^0 \to \mathbf{B}\bar{\mathbf{B}}' D^{*+})/\mathcal{B}(\bar{B}^0 \to \mathbf{B}\bar{\mathbf{B}}' D^+) \gtrsim 3$ for class-I (current-type) decays; (ii) for class-II (transition-type) decays, $\mathcal{B}(\bar{B}^0 \to \mathbf{B}\bar{\mathbf{B}}' D^{*0})/\mathcal{B}(\bar{B}^0 \to \mathbf{B}\bar{\mathbf{B}}' D^0) \gtrsim 1$ or $\lesssim 1$, depending on the modes under consideration; and (iii) $\mathcal{B}(B^- \to \mathbf{B}\bar{\mathbf{B}}' D^{*0})/\mathcal{B}(B^- \to \mathbf{B}\bar{\mathbf{B}}' D^0) \gtrsim 3$ for hybrid-type decays except for $\mathbf{B}\bar{\mathbf{B}}' = \Xi^-\bar{\Sigma}^0$, $\Xi^0\bar{\Sigma}^+$, and $\Xi^-\bar{\Lambda}$.

To get insight into the above-mentioned first feature, we take $\mathbf{B}\bar{\mathbf{B}}' = n\bar{p}$ as an example. In the dibaryon rest frame we have $\vec{p}_{\mathbf{B}} = -\vec{p}_{\bar{\mathbf{B}}'}$ and $\vec{p}_B = \vec{p}_{D^{(*)}}$. As the branching fractions are dominated by the threshold effect, we can consider the dibaryon in the nonrelativistic limit where $q_\mu = (p_{\mathbf{B}} + p_{\bar{\mathbf{B}}'})_\mu \approx (2m, 0, 0, 0)$ with $E_{\mathbf{B}}(E_{\bar{\mathbf{B}}'}) \approx m$. Then we find

$$\begin{aligned}
\langle \mathbf{B}\bar{\mathbf{B}}' | \bar{q}_1 \gamma_\mu q_2 | 0 \rangle &\approx -2m(0, 1, \pm i, 0) F_1(4m^2), \\
\langle \mathbf{B}\bar{\mathbf{B}}' | \bar{q}_1 \gamma_\mu \gamma_5 q_2 | 0 \rangle &\approx -2m(1, 0, 0, 0)(g_A(4m^2) \\
&\quad + h_A(4m^2)).
\end{aligned} \tag{27}$$

However, the vacuum to $\mathbf{B}\bar{\mathbf{B}}'$ transition induced by the axial-vector current is suppressed as $h_A(4m^2) \approx -g_A(4m^2)$ followed from Eq. (17). From Eqs. (6) and (14) we see that the decay amplitudes $\bar{B}^0 \to n\bar{p}D^{*+}$ and $\bar{B}^0 \to n\bar{p}D^{*0}$ are governed by the form factors $A_1^{BD^*}$ and F_1^{BD}, respectively. Thus the ratio of these amplitudes is given by

$$\frac{\mathcal{A}(\bar{B}^0 \to n\bar{p}D^{*+})}{\mathcal{A}(\bar{B}^0 \to n\bar{p}D^{+})} \approx \frac{m_B + m_{D^*}}{2m_D} \frac{E_{D^*}}{|\vec{p}_D|}, \tag{28}$$

where we have assumed $a_1^{D^*} = a_1^D$ and taken the approximation of $A_1^{BD^*}/F_1^{BD} \simeq 1$ inferred from Table III. Note that since numerically $(m_B + m_{D^*})^2/(4m_D^2) \sim 3$, we thus have $\mathcal{B}(\bar{B}^0 \to n\bar{p}D^{*+})/\mathcal{B}(\bar{B}^0 \to n\bar{p}D^+) \sim 3$ as the ratio $E_{D^*}/|\vec{p}_D|$ is close to 1. For heavy dibaryons, the ratio $\mathcal{B}(\bar{B}^0 \to \mathbf{B}\bar{\mathbf{B}}' D^{*+})/\mathcal{B}(\bar{B}^0 \to \mathbf{B}\bar{\mathbf{B}}' D^+)$ will become larger, for example, $\mathcal{B}(\bar{B}^0 \to \Xi^-\bar{\Xi}^0 D^{*+})/\mathcal{B}(\bar{B}^0 \to \Xi^-\bar{\Xi}^0 D^+) \sim [(m_B + m_{D^*})E_{D^*}/(2m_D|\vec{p}_D|)]^2 = 4.2$. Among the current-type modes listed in Table IV, we notice that $\bar{B}^0 \to \Lambda\bar{p}D^{*+}$ and $\bar{B}^0 \to \Sigma^0\bar{p}D^{*+}$ have the largest rates

$$\begin{aligned}
\mathcal{B}(\bar{B}^0 \to \Lambda\bar{p}D^{*+}) &= 1.2 \times 10^{-5}, \\
\mathcal{B}(\bar{B}^0 \to \Sigma^0\bar{p}D^{*+}) &= 1.1 \times 10^{-5}.
\end{aligned} \tag{29}$$

A search of them at B factories is strongly encouraged. Note that the experimental studies of the angular distributions in the charmful cases may help solve the angular correlation puzzle in $\bar{B}^0 \to \Lambda\bar{p}\pi^+$ owing to the same $0 \to \mathbf{B}\bar{\mathbf{B}}'$ transition form factors appearing in both cases.

Similar to the charmless cases of $B^- \to p\bar{p}K^{*-}$ and $\bar{B}^0 \to p\bar{p}K^{*0}$ [40,41], the decay rate of the class-II decay $\bar{B}^0 \to p\bar{p}D^{*0}$ is dominated by the form factors f_2 and g_2 defined in Eq. (18). Contrary to the D^* production, the decay $\bar{B}^0 \to p\bar{p}D^0$ receives contributions from the form factors f_3 and g_3 accompanied by m_D^2. Considering the dibaryon in the nonrelativistic limit, the ratio of the amplitudes reads

$$\frac{\mathcal{A}(\bar{B}^0 \to p\bar{p}D^{*0})}{\mathcal{A}(\bar{B}^0 \to p\bar{p}D^0)} = \frac{g_1|\vec{p}_{D^*}| + f_1 E_{D^*}\sin\theta_p + 2g_2 m_{D^*}(\sqrt{2}|\vec{p}_{D^*}|\cos\theta_p + m_{D^*}\sin\theta_p)}{g_1 E_D + f_1|\vec{p}_D|\sin\theta_p + g_3 m_D^2}. \tag{30}$$

This ratio is expected to be of the order of unity since g_1 and f_1 terms in Eq. (30) are of the same order, and so are g_2 and g_3 terms.

Class-III decays will have rates larger than class-I and class-II ones if the relative phase between a_1 and a_2 is small so that their interference is constructive. To proceed, we use $a_1^{D^{(*)}}$ and $a_2^{D^{(*)}}$ extracted from $\bar{B}^0 \to n\bar{p}D^{*+}$ and $\bar{B}^0 \to p\bar{p}D^{*0}$, respectively, and neglect their phases. The results are summarized in Table VI. For example,

$$\begin{aligned}\mathcal{B}(B^- \to \Lambda\bar{p}D^0) &= 1.1 \times 10^{-5},\\ \mathcal{B}(B^- \to \Lambda\bar{p}D^{*0}) &= 3.2 \times 10^{-5},\end{aligned} \tag{31}$$

which are larger than $\mathcal{B}(\bar{B}^0 \to \Lambda\bar{p}D^+)$ and $\mathcal{B}(\bar{B}^0 \to \Lambda\bar{p}D^{*+})$, respectively, by a factor of 3. It will be interesting to test this experimentally.

As seen in Tables VII and VIII, for the baryonic decays with a J/ψ in the final state, our prediction $\mathcal{B}(B^- \to \Sigma^0\bar{p}J/\Psi) = 1.3 \times 10^{-7}$ is consistent with the Belle limit, 1.1×10^{-5}, and $\mathcal{B}(\bar{B}^0 \to p\bar{p}J/\Psi) = 1.2 \times 10^{-6}$ is in accordance with the *BABAR* limit but slightly higher than the upper bound set by Belle [see Eq. (2) for the data]. We also find that the decays $B^- \to (\Xi^-\bar{\Sigma}^0, \Xi^-\Lambda, \Xi^0\bar{\Sigma}^+, \Xi^0\bar{\Sigma}^0, \Xi^-\bar{\Sigma}^-)J/\Psi$ have branching fractions of order 10^{-5}.

As for the threshold peaking effect, while it manifests in the decay $\bar{B}^0 \to p\bar{p}D^{(*)0}$ the data clearly do not show the threshold behavior in $B^- \to \Lambda\bar{p}J/\Psi$ (see Fig. 2). This can be understood as follows. In the latter decay, the invariant mass $m_{\Lambda\bar{p}}$ ranges from 2.05 to 2.18 GeV, which is very narrow compared to the $m_{p\bar{p}}$ range in the $\bar{B}^0 \to p\bar{p}D^{(*)0}$ decay. Consequently, the invariant mass distribution of $d\Gamma/dm_{\Lambda\bar{p}}$ is governed by the shape of the phase space due to the relative flat $1/t^3$ dependence within the small allowed $m_{\Lambda\bar{p}}$ region.

Finally we notice that since the transition-type decays share the same $B \to \mathbf{B}\bar{\mathbf{B}}'$ transition form factors with $B^- \to p\bar{p}\pi^-$, $B^- \to p\bar{p}K^-$, and $B^- \to \Lambda\bar{p}\gamma$, measurements for distributions/Daltiz plot asymmetries are useful for solving the puzzle of the angular distributions in the charmless baryonic B decays. For example, a recent Dalitz plot analysis of $\bar{B}^0 \to p\bar{p}D^{(*)0}$ by *BABAR* [24] shows a possible indication that

$$A_\theta(\bar{B}^0 \to p\bar{p}D^0) > A_\theta(\bar{B}^0 \to p\bar{p}D^{*0}) > 0, \tag{32}$$

where the angular asymmetry A_θ is defined by [51]

$$A_\theta \equiv \frac{\int_0^{+1} \frac{d\Gamma}{d\cos\theta} d\cos\theta - \int_{-1}^{0} \frac{d\Gamma}{d\cos\theta} d\cos\theta}{\int_0^{+1} \frac{d\Gamma}{d\cos\theta} d\cos\theta + \int_{-1}^{0} \frac{d\Gamma}{d\cos\theta} d\cos\theta}, \tag{33}$$

which is equal to $(N_+ - N_-)/(N_+ + N_-)$, where $N_\pm$ are the events with $\cos\theta > 0$ and $\cos\theta < 0$, respectively. The above relation (32) can be explained in our model: in the previous study in [41,51] we have predicted that $|A_\theta(B^- \to p\bar{p}K^-)| > |A_\theta(B^- \to p\bar{p}K^{*-})|$ and $|A_\theta(B^- \to p\bar{p}\pi^-)| > |A_\theta(B^- \to p\bar{p}\rho^-)|$. In our model, the charmful baryonic B decays have similar features of the angular distribution as the charmless cases. Therefore, we believe that our model together with the choice of the form factors will help understand the angular correlations observed in these decays.

VI. CONCLUSIONS

Within the framework of the generalized factorization approach, we have studied three different types of charmful three-body baryonic B decays with $D^{(*)}$ or J/Ψ in the final state. We found that the $B \to \mathbf{B}\bar{\mathbf{B}}'$ form factors extracted mainly from the data of charmless baryonic B decays can be used to explain the measured rates of $\bar{B}^0 \to p\bar{p}D^{(*)0}$ and $B^- \to \Lambda\bar{p}J/\Psi$. The branching fractions of $\bar{B}^0 \to \Lambda\bar{p}D^{*+}$, $\bar{B}^0 \to \Sigma^0\bar{p}D^{*+}$, $B^- \to \Lambda\bar{p}D^0$, and $B^- \to \Lambda\bar{p}D^{*0}$ are predicted to be $(1.2, 1.1, 1.1, 3.9) \times 10^{-5}$, respectively. They are readily accessible to B factories.

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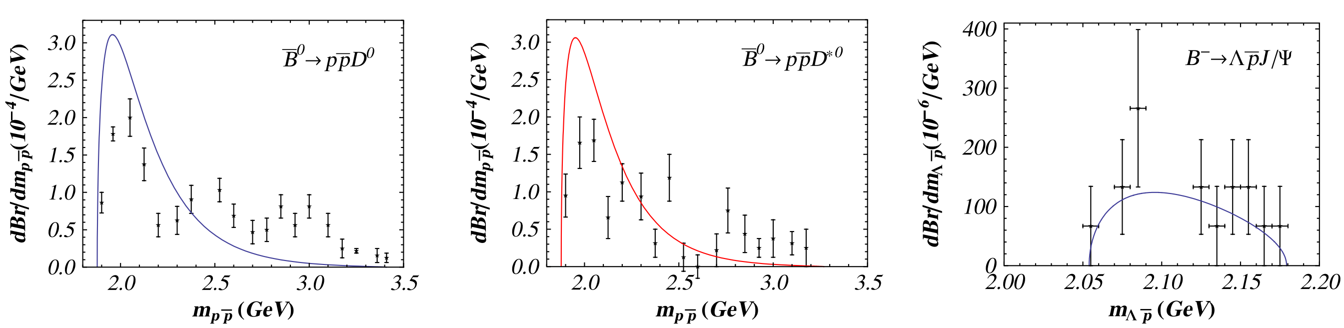

FIG. 2 (color online). Dibaryon invariant mass distributions for $\bar{B}^0 \to p\bar{p}D^0$, $\bar{B}^0 \to p\bar{p}D^{*0}$, and $B^- \to \Lambda\bar{p}J/\Psi$, respectively. Experimental data are taken from [24,27].

APPENDIX: RELATIONS FOR $B \to \mathbf{B}\bar{\mathbf{B}}'$ TRANSITION FORM FACTORS

Form factors of $\langle \mathbf{B}\bar{\mathbf{B}}'|\bar{q}\gamma_\mu b|B\rangle$ and $\langle \mathbf{B}\bar{\mathbf{B}}'|\bar{q}\gamma_\mu\gamma_5 b|B\rangle$ can be related to each other by the SU(3) flavor and SU(2) spin symmetries in the large t limit ($t \to \infty$). As the currents $V_\mu(A_\mu) = \bar{q}\gamma_\mu(\gamma_5)b$ are used to create the dibaryon and annihilate the B meson with $|B\rangle \sim |\bar{b}\gamma_5 q'|0\rangle$, we have

$$\begin{aligned}\langle \mathbf{B}\bar{\mathbf{B}}'|V_\mu - A_\mu|B\rangle &= \langle \mathbf{B}\bar{\mathbf{B}}'|2\bar{q}_L\gamma_\mu b_L(\bar{b}_L q'_R - \bar{b}_R q'_L)|0\rangle \\ &= \langle \mathbf{B}\bar{\mathbf{B}}'|2\bar{q}_L\gamma_\mu(\not{p}_b + m_b)q'_R \\ &\quad - 2\bar{q}_L\gamma_\mu(\not{p}_b + m_b)q'_L|0\rangle \\ &= \langle \mathbf{B}\bar{\mathbf{B}}'| - 2m_b\bar{q}_L\gamma_\mu q'_L \\ &\quad + 2\bar{q}_L\gamma_\mu\not{p}_b q'_R|0\rangle \\ &= \langle \mathbf{B}\bar{\mathbf{B}}'|J_\mu|0\rangle + \langle \mathbf{B}\bar{\mathbf{B}}'|J'_\mu|0\rangle, \qquad \text{(A1)}\end{aligned}$$

where $J_\mu = -2m_b\bar{q}_L\gamma_\mu q'_L$ and $J'_\mu = 2\bar{q}_L\gamma_\mu\not{p}_b q'_R$. Note that J_μ (J'_μ) is the source of the chiral-even (chiral-odd) current. In terms of the crossing symmetry, the antibaryon in the final state can be transformed as the baryon in the initial state. Then we can follow Refs. [37,38,48] to parametrize the $\mathbf{B}' \to \mathbf{B}$ transition form factors as

$$\begin{aligned}\langle \mathbf{B}|J_\mu|\mathbf{B}'\rangle &= 2im_b\bar{u}\gamma_\mu\left[\frac{1+\gamma_5}{2}D^+ + \frac{1-\gamma_5}{2}D^-\right]u, \\ \langle \mathbf{B}|J'_\mu|\mathbf{B}'\rangle &= 2i\bar{u}\gamma_\mu\not{p}_b\left[\frac{1+\gamma_5}{2}D'^+ + \frac{1-\gamma_5}{2}D'^-\right]u. \qquad \text{(A2)}\end{aligned}$$

The above chiral-even and chiral-odd form factors $D^\pm$ and $D'^\pm$ are given by

$$D^\pm = e^\pm_\parallel D_\parallel + e^\pm_{\overline{\parallel}} D_{\overline{\parallel}}, \qquad D'^\pm = e^\pm_\parallel D'_\parallel + e^\pm_{\overline{\parallel}} D'_{\overline{\parallel}}, \quad \text{(A3)}$$

where $D_{\parallel(\overline{\parallel})}$, $D'_{\parallel(\overline{\parallel})}$ are the amplitudes in which the currents interact with a valence quark with helicity parallel (antiparallel) to the helicity of the baryon. The parameters $e^\pm_\parallel$ and $e^\pm_{\overline{\parallel}}$ are defined by

$$\begin{aligned}e^\pm_\parallel &= \sum_{i<j}\langle \mathbf{B}, h|Q(i) + Q(j)|\mathbf{B}', h'\rangle, \\ e^\pm_{\overline{\parallel}} &= \sum_k \langle \mathbf{B}, h|Q(k)|\mathbf{B}', h'\rangle, \qquad \text{(A4)}\end{aligned}$$

where the charge $Q = J_0$ with $i, j, k = 1, 2, 3$ numbering the valence quarks in the baryon to be acted. Therefore, $e^\pm_{\parallel(\overline{\parallel})}$ is the sum of the electroweak charges carried by valence quarks in the baryon with helicities parallel (antiparallel) to the baryon's helicity. Besides, the superscript $+(-)$ denotes $h' = \uparrow(\downarrow)$, the helicities of the initial baryon.

Take $\bar{B}^0 \to p\bar{p}$ as an example. With the helicity state of the proton given by

$$\begin{aligned}|p,\uparrow\rangle &= |p,\uparrow\downarrow\uparrow\rangle + |p,\uparrow\uparrow\downarrow\rangle + |p,\downarrow\uparrow\uparrow\rangle \\ &= \frac{1}{\sqrt{3}}\left\{\frac{1}{\sqrt{6}}[2u_\uparrow d_\downarrow u_\uparrow - u_\uparrow u_\downarrow d_\uparrow - d_\uparrow u_\downarrow u_\uparrow]\right. \\ &\quad + \frac{1}{\sqrt{6}}[2u_\uparrow u_\uparrow d_\downarrow - u_\uparrow d_\uparrow u_\downarrow - d_\uparrow u_\uparrow u_\downarrow] \\ &\quad \left. + \frac{1}{\sqrt{6}}[2d_\downarrow u_\uparrow u_\uparrow - u_\downarrow d_\uparrow u_\uparrow - u_\downarrow u_\uparrow d_\uparrow]\right\}, \\ |p,\downarrow\rangle &= (|p,\uparrow\rangle \quad \text{with} \quad \uparrow\leftrightarrow\downarrow), \qquad \text{(A5)}\end{aligned}$$

we obtain

$$\begin{aligned}e^-_\parallel &= \sum_{i<j}\langle p,\downarrow|Q(i) + Q(j)|p,\downarrow\rangle = \frac{1}{3}, \\ e^+_\parallel &= \sum_{i<j}\langle p,\uparrow|Q(i) + Q(j)|p,\uparrow\rangle = 0, \\ e^-_{\overline{\parallel}} &= \sum_k\langle p,\downarrow|Q(k)|p,\downarrow\rangle = 0, \\ e^+_{\overline{\parallel}} &= \sum_k\langle p,\uparrow|Q(k)|p,\uparrow\rangle = \frac{2}{3}, \qquad \text{(A6)}\end{aligned}$$

for the chiral-even J_μ, and

$$\begin{aligned}e^-_\parallel &= \sum_{i<j}\langle p,\downarrow|Q(i) + Q(j)|p,\downarrow\rangle = 0, \\ e^+_\parallel &= \sum_{i<j}\langle p,\uparrow|Q(i) + Q(j)|p,\uparrow\rangle = \frac{-1}{3}, \\ e^-_{\overline{\parallel}} &= \sum_k\langle p,\downarrow|Q(k)|p,\downarrow\rangle = 0, \\ e^+_{\overline{\parallel}} &= \sum_k\langle p,\uparrow|Q(k)|p,\uparrow\rangle = 0, \qquad \text{(A7)}\end{aligned}$$

for the chiral-odd J'_μ. When the sums of the electroweak charges $e^\pm_{\parallel(\overline{\parallel})}$ derived in Eqs. (A6) and (A7) are put back into Eq. (A2), $\langle \mathbf{B}|J_\mu + J'_\mu|\mathbf{B}'\rangle$ is given as

$$\begin{aligned}\langle \mathbf{B}|J_\mu + J'_\mu|\mathbf{B}'\rangle &= i\bar{u}\{m_b[\tfrac{1}{3}D_\parallel - \tfrac{2}{3}D_{\overline{\parallel}}]\gamma_\mu \\ &\quad + [-\tfrac{1}{3}D'_\parallel]\gamma_\mu\not{p}_b\}\gamma_5 u \\ &\quad + i\bar{u}\{m_b[\tfrac{1}{3}D_\parallel + \tfrac{2}{3}D_{\overline{\parallel}}]\gamma_\mu \\ &\quad + [\tfrac{1}{3}D'_\parallel]\gamma_\mu\not{p}_b\}u. \qquad \text{(A8)}\end{aligned}$$

This gives the form factor relations for $\langle \mathbf{B}\bar{\mathbf{B}}'|\bar{q}\gamma_\mu b|B\rangle$ and $\langle \mathbf{B}\bar{\mathbf{B}}'|\bar{q}\gamma_\mu\gamma_5 b|B\rangle$, which is given by

$$\begin{aligned}g_1 &= \tfrac{1}{3}D_\parallel - \tfrac{2}{3}D_{\overline{\parallel}}, \qquad f_1 = \tfrac{1}{3}D_\parallel + \tfrac{2}{3}D_{\overline{\parallel}}, \\ g_i &= -f_i = -\tfrac{1}{3}D^i_\parallel, \qquad (i = 2, \ldots, 5). \qquad \text{(A9)}\end{aligned}$$

By the same token, we find the relation of $g_i = -f_i$ holds in all $B \to \mathbf{B}\bar{\mathbf{B}}'$ transition form factors.

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