Violation of partial conservation of the axial-vector current and two-body baryonic B and D_s decays

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We study the two-body baryonic B and D_s decays based on the annihilation mechanism without the partial conservation of the axial-vector current (PCAC) at the GeV scale. We demonstrate that the contributions of $B^- \to \Lambda\bar{p}$, $B^- \to \Sigma^0\bar{p}$, and $\bar{B}^0_s \to \Lambda\bar{\Lambda}$ are mainly from the scalar and pseudoscalar currents with their branching ratios predicted to be around $(3.5, 5.3, 5.3) \times 10^{-8}$, respectively, exactly the sizes of $\mathcal{B}(B \to \mathbf{B}\bar{\mathbf{B}}')$ established by the data. We also apply the annihilation mechanism to all of the charmless two-body baryonic B and D_s decays. In particular, we can explain $\mathcal{B}(\bar{B}^0_{(s)} \to p\bar{p})$ of order 10^{-8} and $\mathcal{B}(D^+_s \to p\bar{n})$ of order 10^{-3}, which are from the axial-vector currents. In addition, the branching ratios of $\bar{B}^0 \to \Lambda\bar{\Lambda}$, $B^- \to n\bar{p}$, and $B^- \to \Sigma^-\bar{\Sigma}^0$ are predicted to be $(0.3, 3.2, 9.6) \times 10^{-8}$, which can be measured by LHCb and viewed as tests for the violation of the partial conservation of the axial-vector current at the GeV scale.

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I. INTRODUCTION

For the abundantly observed three-body baryonic B decays ($B \to \mathbf{B}\bar{\mathbf{B}}'M$), the theoretical approach for the systematic study has been established [1–6]. It leads to the theoretical predictions, among which at least five decay modes [7,8] are observed to agree with the data [9]. On the other hand, the two-body baryonic B decays ($B \to \mathbf{B}\bar{\mathbf{B}}'$) are poorly understood due to the smaller branching ratios, causing a much later observation than $B \to \mathbf{B}\bar{\mathbf{B}}'M$. Recently, the LHCb collaboration has presented the first observations of the charmless $B \to \mathbf{B}\bar{\mathbf{B}}'$ decays [10], given by

$$\mathcal{B}(\bar{B}^0 \to p\bar{p}) = (1.47^{+0.62+0.35}_{-0.51-0.14}) \times 10^{-8},$$
$$\mathcal{B}(\bar{B}^0_s \to p\bar{p}) = (2.84^{+2.03+0.85}_{-1.68-0.18}) \times 10^{-8}, \qquad (1)$$

with the statistical significances to be 3.3σ and 1.9σ, respectively.

Based on the factorization, when the B meson annihilates with the momentum transfer q, the amplitudes $\mathcal{A}(\bar{B}^0_{(s)} \to p\bar{p})$ can be decomposed as $q^\mu\langle p\bar{p}|A_\mu|0\rangle$, where the matrix element is for the proton pair production and A_μ is the axial-vector current. From the hypothesis of the partial conservation of the axial-vector current (PCAC) [11] at the GeV scale, $q^\mu A_\mu$ is proportional to m^2_π, which leads to $\mathcal{A}(\bar{B}^0_{(s)} \to p\bar{p}) \simeq 0$. This is the reason why the nonfactorizable effects were believed to dominate the branching ratios in Eq. (1) [12].[1] However, since the predictions from these models differ from each other, and commonly exceed the data, a reliable theoretical approach has not been established yet.

In this work, we would propose a new method without the use of the PCAC. In fact, the smallness of the previous estimations is not caused by the annihilation mechanism [13] but the assumption of PCAC. Moreover, this assumption has never been tested at the GeV scale. For example, $\mathcal{B}(B^- \to \Lambda\bar{p})$ and $\mathcal{B}(\bar{B}^0_s \to \Lambda\bar{\Lambda})$ are found to have the amplitudes decomposed as $(m^2_B/m_b)\langle p\bar{p}|S+P|0\rangle$ with $S(P)$ the (pseudo)scalar current, which has no connection to the PCAC. Since they can be estimated to be of order 10^{-8}, exactly the order of the magnitude of $\mathcal{B}(B \to \mathbf{B}\bar{\mathbf{B}}')$ measured by the experiments, the annihilation mechanism can be justified. If the axial-vector current is asymptotically conserved, the result of $\mathcal{B}(D^+_s \to p\bar{n}) = (0.4^{+1.1}_{-0.3}) \times 10^{-6}$ in Ref. [14] would yield $\mathcal{B}(D^+_s \to p\bar{n})/\mathcal{B}(D^+_s \to \tau\bar{\nu}_\tau) \simeq 10^{-5}$, which was indeed suggested as the test of the PCAC at the GeV scale [13]. Nonetheless, with $\mathcal{B}(D^+_s \to p\bar{n}) = (1.30 \pm 0.36^{+0.12}_{-0.16}) \times 10^{-3}$ measured by the CLEO Collaboration [15], one obtains that $\mathcal{B}(D^+_s \to p\bar{n})/\mathcal{B}(D^+_s \to \tau\bar{\nu}_\tau) \simeq 0.02$, which is too large and can be viewed as a countercase of the PCAC [16].

In this paper, we apply the annihilation mechanism to the two-body baryonic B decays, provided that the axial-vector current is not asymptotically conserved. By modifying the timelike baryonic form factors via the axial-vector current without respect to the PCAC, we can explain $\mathcal{B}(\bar{B}^0_{(s)} \to p\bar{p})$ as well as $\mathcal{B}(D^+_s \to p\bar{n})$. We shall also predict $\mathcal{B}(B^- \to \Lambda(\Sigma^0)\bar{p})$ and $\mathcal{B}(\bar{B}^0_s \to \Lambda\bar{\Lambda})$ in terms of the timelike baryonic form factors via the scalar and pseudoscalar currents.

The paper is organized as follows. In Sec. II, we present the formalism of the two-body baryonic B and D_s decays.

[1]For the review of the various models, please consult Ref. [12] and the references therein.

In Sec. III, we proceed with our numerical analysis. Section IV contains our discussions and conclusions.

II. FORMALISM

The nonleptonic B and D decays in the factorization hypothesis are in analogy with the semileptonic cases like $\mathcal{A}(B \to \pi e \bar{\nu}_e) \propto \langle \pi | u \gamma^\mu (1-\gamma_5) b | B \rangle \bar{e} \gamma_\mu (1-\gamma_5) \nu_e$ to have the amplitudes with an additional matrix element in the form of $\langle X_2 | J^{2(\mu)} | 0 \rangle \langle X_1 | J^1_{(\mu)} | B \rangle$, where $J^{1,2}_{(\mu)}$ are the quark currents, and $X_{1,2}$ can be multihadron states [17,18]. Although the derivation may not be analytically satisfactory, the factorization approximation can still be justified by theoretically reproducing the data and predicting not-yet-observed decay modes to be approved by the later measurements in the two-body and three-body mesonic B decays as well as the three-body baryonic B decays [8,19–21].

Like the measured $\bar{B}^0_{(s)} \to p\bar{p}$ and $D_s^+ \to p\bar{n}$ with the decaying processes depicted in Fig. 1, in the two-body baryonic B and D_s decays, the factorizable amplitudes are known to depend on the annihilation mechanism [13,16], where B and D_s annihilate, followed by the baryon pair production. Thus, the amplitudes can have two types, $\mathcal{A}_1$ and $\mathcal{A}_2$, which consist of (axial)vectors and (pseudo)scalar quark currents, respectively. For example, the amplitudes of $\bar{B}^0 \to (p\bar{p}, \Lambda\bar{\Lambda})$, $B^- \to (n\bar{p}, \Sigma^-\bar{\Sigma}^0)$, and $D_s^+ \to p\bar{n}$ are of the first type, given by [13,14,16]

$$\begin{aligned}
\mathcal{A}_1(\bar{B}^0 \to \mathbf{B_1}\bar{\mathbf{B}}'_1) &= \frac{G_F}{\sqrt{2}} V_{ub} V^*_{ud} a_2 \\
&\quad \times \langle \mathbf{B_1}\bar{\mathbf{B}}'_1 | (\bar{u}u)_{V-A} | 0 \rangle \langle 0 | (\bar{d}b)_{V-A} | \bar{B}^0 \rangle, \\
\mathcal{A}_1(B^- \to \mathbf{B_2}\bar{\mathbf{B}}'_2) &= \frac{G_F}{\sqrt{2}} V_{ub} V^*_{ud} a_1 \\
&\quad \times \langle \mathbf{B_2}\bar{\mathbf{B}}'_2 | (\bar{d}u)_{V-A} | 0 \rangle \langle 0 | (\bar{u}b)_{V-A} | B^- \rangle, \\
\mathcal{A}_1(D_s^+ \to p\bar{n}) &= \frac{G_F}{\sqrt{2}} V_{cs} V^*_{ud} a_1 \\
&\quad \times \langle p\bar{n} | (\bar{u}d)_{V-A} | 0 \rangle \langle 0 | (\bar{s}c)_{V-A} | D_s^+ \rangle, \qquad (2)
\end{aligned}$$

where $\mathbf{B_1}\bar{\mathbf{B}}'_1 = p\bar{p}$ or $\Lambda\bar{\Lambda}$, $\mathbf{B_2}\bar{\mathbf{B}}'_2 = n\bar{p}$ or $\Sigma^-\bar{\Sigma}^0$, $(\bar{q}_1 q_2)_{V-A}$ denotes $\bar{q}_1 \gamma_\mu (1-\gamma_5) q_2$, G_F is the Fermi constant, a_i are the effective Wilson coefficients, and $V_{q_1 q_2}$ are the Cabibbo–Kobayashi–Maskawa (CKM) matrix elements. The amplitudes of $\bar{B}^0_s \to (p\bar{p}, \Lambda\bar{\Lambda})$ and $B^- \to (\Lambda\bar{p}, \Sigma^0\bar{p})$ are more complicated, written as

$$\begin{aligned}
\mathcal{A}(\bar{B}^0_s \to \mathbf{B_1}\bar{\mathbf{B}}'_1) &= \mathcal{A}_1(\bar{B}^0_s \to \mathbf{B_1}\bar{\mathbf{B}}'_1) + \mathcal{A}_2(\bar{B}^0_s \to \mathbf{B_1}\bar{\mathbf{B}}'_1), \\
\mathcal{A}(B^- \to \mathbf{B_2}\bar{\mathbf{B}}'_2) &= \mathcal{A}_1(B^- \to \mathbf{B_2}\bar{\mathbf{B}}'_2) + \mathcal{A}_2(B^- \to \mathbf{B_2}\bar{\mathbf{B}}'_2), \qquad (3)
\end{aligned}$$

where

$$\begin{aligned}
\mathcal{A}_1(\bar{B}^0_s \to \mathbf{B_1}\bar{\mathbf{B}}'_1) &= \frac{G_F}{\sqrt{2}} \Big\{ V_{ub} V^*_{us} a_2 \langle \mathbf{B_1}\bar{\mathbf{B}}'_1 | (\bar{u}u)_{V-A} | 0 \rangle - V_{tb} V^*_{ts} \Big[a_3 \langle \mathbf{B_1}\bar{\mathbf{B}}'_1 | (\bar{u}u + \bar{d}d + \bar{s}s)_{V-A} | 0 \rangle \\
&\quad + a_4 \langle \mathbf{B_1}\bar{\mathbf{B}}'_1 | (\bar{s}s)_{V-A} | 0 \rangle + a_5 \langle \mathbf{B_1}\bar{\mathbf{B}}'_1 | (\bar{u}u + \bar{d}d + \bar{s}s)_{V+A} | 0 \rangle \\
&\quad + \frac{a_9}{2} \langle \mathbf{B_1}\bar{\mathbf{B}}'_1 | (2\bar{u}u - \bar{d}d - \bar{s}s)_{V-A} | 0 \rangle \Big] \Big\} \langle 0 | (\bar{s}b)_{V-A} | \bar{B}^0_s \rangle, \\
\mathcal{A}_1(B^- \to \mathbf{B_2}\bar{\mathbf{B}}'_2) &= \frac{G_F}{\sqrt{2}} (V_{ub} V^*_{us} a_1 - V_{tb} V^*_{ts} a_4) \langle \mathbf{B_2}\bar{\mathbf{B}}'_2 | (\bar{s}u)_{V-A} | 0 \rangle \langle 0 | (\bar{u}b)_{V-A} | B^- \rangle, \qquad (4)
\end{aligned}$$

and

$$\begin{aligned}
\mathcal{A}_2(\bar{B}^0_s \to \mathbf{B_1}\bar{\mathbf{B}}'_1) &= \frac{G_F}{\sqrt{2}} V_{tb} V^*_{ts} 2a_6 \\
&\quad \times \langle \mathbf{B_1}\bar{\mathbf{B}}'_1 | (\bar{s}s)_{S+P} | 0 \rangle \langle 0 | (\bar{s}b)_{S-P} | \bar{B}^0_s \rangle, \\
\mathcal{A}_2(B^- \to \mathbf{B_2}\bar{\mathbf{B}}'_2) &= \frac{G_F}{\sqrt{2}} V_{tb} V^*_{ts} 2a_6 \\
&\quad \times \langle \mathbf{B_2}\bar{\mathbf{B}}'_2 | (\bar{s}u)_{S+P} | 0 \rangle \langle 0 | (\bar{u}b)_{S-P} | B^- \rangle, \qquad (5)
\end{aligned}$$

with $\mathbf{B_1}\bar{\mathbf{B}}'_1 = p\bar{p}$ or $\Lambda\bar{\Lambda}$, $\mathbf{B_2}\bar{\mathbf{B}}'_2 = \Lambda\bar{p}$ or $\Sigma^0\bar{p}$ and $(\bar{q}_1 q_2)_{S\pm P}$ representing $\bar{q}_1(1\pm\gamma_5)q_2$. For the coefficients a_i in Eqs. (2)–(5), we use the same inputs as those in $B \to \mathbf{B}\bar{\mathbf{B}}'M$ [7,8], where $a_i = c_i^{\rm eff} + c_{i\pm1}^{\rm eff}/N_c$ with the color number N_c for $i =$ odd (even) in terms of the effective Wilson coefficients $c_i^{\rm eff}$, defined in Refs. [19,20]. Note that N_c is floating between 2 and ∞ in the generalized factorization for the correction of the nonfactorizable effects. In Eqs. (3)–(5), the matrix element for the annihilation of the pseudoscalar meson is defined by

$$\langle 0 | \bar{q}_1 \gamma_\mu \gamma_5 q_2 | P \rangle = i f_P q_\mu, \qquad (6)$$

with f_P the decay constant, from which we can obtain $\langle 0 | \bar{q}_1 \gamma_5 q_2 | P \rangle$ by using the equation of motion: $-i\partial^\mu(\bar{q}_1\gamma_\mu q_2) = (m_{q_1} - m_{q_2})\bar{q}_1 q_2$ and $-i\partial^\mu(\bar{q}_1\gamma_\mu\gamma_5 q_2) = (m_{q_1} + m_{q_2})\bar{q}_1\gamma_5 q_2$. For the dibaryon production, the matrix elements read

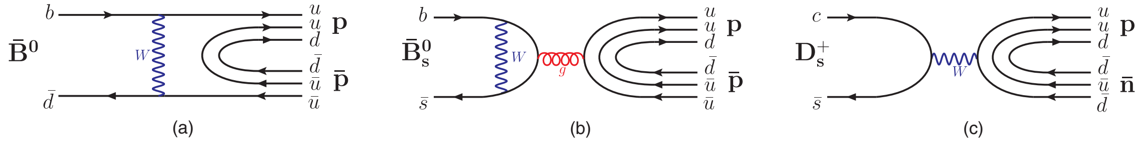

FIG. 1 (color online). The two-body baryonic decays of (a) $\bar{B}^0 \to p\bar{p}$, (b) $\bar{B}^0_s \to p\bar{p}$, and (c) $D^+_s \to p\bar{n}$.

$$\langle \mathbf{B}\bar{\mathbf{B}}'|\bar{q}_1\gamma_\mu q_2|0\rangle = \bar{u}\left\{F_1\gamma_\mu + \frac{F_2}{m_{\mathbf{B}}+m_{\bar{\mathbf{B}}'}}i\sigma_{\mu\nu}q_\mu\right\}v,$$
$$\langle \mathbf{B}\bar{\mathbf{B}}'|\bar{q}_1\gamma_\mu\gamma_5 q_2|0\rangle = \bar{u}\left\{g_A\gamma_\mu + \frac{h_A}{m_{\mathbf{B}}+m_{\bar{\mathbf{B}}'}}q_\mu\right\}\gamma_5 v,$$
$$\langle \mathbf{B}\bar{\mathbf{B}}'|\bar{q}_1 q_2|0\rangle = f_S\bar{u}v,$$
$$\langle \mathbf{B}\bar{\mathbf{B}}'|\bar{q}_1\gamma_5 q_2|0\rangle = g_P\bar{u}\gamma_5 v, \tag{7}$$

with $u(v)$ as the (anti)baryon spinor, where $F_{1,2}$, g_A, h_A, f_S, and g_P are the timelike baryonic form factors. The amplitudes $\mathcal{A}_1$ and $\mathcal{A}_2$ now can be reduced as

$$\mathcal{A}_1 \propto \frac{1}{(m_{\mathbf{B}}+m_{\bar{\mathbf{B}}'})}\bar{u}[(m_{\mathbf{B}}+m_{\bar{\mathbf{B}}'})^2 g_A + m^2_{B(D_s)}h_A]\gamma_5 v,$$
$$\mathcal{A}_2 \propto \frac{m_B^2}{m_b}\bar{u}(f_S + g_P\gamma_5)v. \tag{8}$$

Note that f_S and g_P are not suppressed by any relations, such that the factorization obviously works for the decay modes with $\mathcal{A}_2$. Besides, the absence of $F_{1,2}$ in $\mathcal{A}_1$ corresponds to the conserved vector current. However, due to the equation of motion, F_1 reappears as a part of f_S in $\mathcal{A}_2$, given by

$$f_S = n_q F_1, \tag{9}$$

with $n_q = (m_{\mathbf{B}} - m_{\mathbf{B}'})/(m_{q_1} - m_{q_2})$, which is fixed to be 1.3 [3,7], presenting 30% of the $SU(3)$ flavor symmetry breaking effect. In perturbative QCD counting rules, the momentum dependences of F_1 and g_A can be written as [22–24]

$$F_1 = \frac{C_{F_1}}{t^2}\left[\ln\left(\frac{t}{\Lambda_0^2}\right)\right]^{-\gamma}, \quad g_A = \frac{C_{g_A}}{t^2}\left[\ln\left(\frac{t}{\Lambda_0^2}\right)\right]^{-\gamma}, \tag{10}$$

with $t \equiv (p_{\mathbf{B}} + p_{\mathbf{B}'})^2$, where $\gamma = 2 + 4/(3\beta) = 2.148$, with β being the QCD β function and $\Lambda_0 = 0.3$ GeV. We note that, as the leading-order expansion, F_1 and g_A ($\propto 1/t^2$) account for two hard gluons, which connect to the valence quarks within the dibaryon. In terms of the PCAC, one obtains the relations of

$$h_A = -\frac{(m_{\mathbf{B}}+m_{\mathbf{B}'})^2}{t-m_M^2}g_A, \qquad g_P = -\frac{m_{\mathbf{B}}+m_{\mathbf{B}'}}{m_{q_1}+m_{q_2}}\frac{m_M^2}{t-m_M^2}g_A, \tag{11}$$

where m_M stands for the meson pole, while g_P is related to g_A from the equation of motion. When h_A in Eq. (11) is used for $B \to \mathbf{B}\bar{\mathbf{B}}'$ with $t = m_B^2 \gg m_M^2$, $\mathcal{B}(\bar{B}^0_{(s)} \to p\bar{p})$ with a suppressed $A_1 \simeq 0$ fails to explain the data by several orders of magnitude. Similarly, $\mathcal{B}(\bar{B}^0 \to \Lambda\bar{p}\pi^+(\rho^+))$ cannot be understood either with g_P in Eq. (11) [2,3]. We hence conclude that h_A and g_P in Eq. (11) from the PCAC at the GeV scale are unsuitable. Recall that F_1 and g_A, where $F_1 = F_1(0)/(1-t/m_V^2)^2$ and $g_A = g_A(0)/(1-t/m_A^2)^2$ [25] with the pole effects for low momentum transfer, have been replaced by Eq. (10) for the decays at the GeV scale. It is reasonable to rewrite h_A and g_P to be

$$h_A = \frac{C_{h_A}}{t^2}, \qquad g_P = f_S, \tag{12}$$

where h_A is inspired by the relation in Eq. (11). For h_A in Eq. (11), since the prefactor, $-(m_{\mathbf{B}}+m_{\mathbf{B}'})^2/t$, arises from the equation of motion, it indicates that both h_A and g_A behave as $1/t^2$. Besides, at the threshold area of $t \simeq (m_{\mathbf{B}}+m_{\mathbf{B}'})^2$, it turns out that $h_A \simeq -g_A$. We regard $h_A = C_{h_A}/t^2$ as the modification of Eq. (11). Consequently, the PCAC is violated; i.e., the axial-vector current is no more asymptotically conserved. As a result of the $SU(3)$ flavor and $SU(2)$ helicity symmetries, $g_P = f_S$ was first derived in Ref. [4], which successfully explained $\mathcal{B}(\bar{B}^0 \to \Lambda\bar{p}\pi^+(\rho^+))$ [4,26].

In Refs. [22–24,26], C_{F_1} and C_{g_A} have been derived carefully to be combined as another set of parameters $C_{\parallel}$ and $C_{\overline{\parallel}}$, which are from the chiral currents. Here, we take the $p\bar{n}$ production for our description. First, due to the crossing symmetry, $\langle p\bar{n}|(\bar{u}d)_{V(A)}|0\rangle$ for the timelike $p\bar{n}$ production and $\langle p|(\bar{u}d)_{V(A)}|n\rangle$ for the spacelike n to p transiton are in fact identical. Therefore, the approach of the pQCD counting rules for the spacelike $\mathbf{B}' \to \mathbf{B}$ transition is useful [24]. We hence combine the vector and axial-vector quark currents, $V_\mu = \bar{u}\gamma_\mu d$ and $A_\mu = \bar{u}\gamma_\mu\gamma_5 d$, to be the right-handed chiral current $J_R^\mu = (V^\mu + A^\mu)/2$, which corresponds to another set of matrix elements for the n to p transition,

$$\langle p_{R+L}|J_R^\mu|n_{R+L}\rangle = \bar{u}\left[\gamma_\mu\frac{1+\gamma_5}{2}G^\uparrow(t) + \gamma_\mu\frac{1-\gamma_5}{2}G^\downarrow(t)\right]u, \tag{13}$$

where the two chiral baryon states $|\mathbf{B}_{R+L}\rangle$ become the two helicity states $|\mathbf{B}_{\uparrow+\downarrow}\rangle \equiv |\mathbf{B}_\uparrow\rangle + |\mathbf{B}_\downarrow\rangle$ in the large t limit. The new set of form factors $G^\uparrow(t)$ and $G^\downarrow(t)$ are defined as

$$G^{\uparrow}(t) = e^{\uparrow}_{\parallel} G_{\parallel}(t) + e^{\uparrow}_{\overline{\parallel}} G_{\overline{\parallel}}(t),$$
$$G^{\downarrow}(t) = e^{\downarrow}_{\parallel} G_{\parallel}(t) + e^{\downarrow}_{\overline{\parallel}} G_{\overline{\parallel}}(t), \qquad (14)$$

where

$$G_{\parallel(\overline{\parallel})}(t) = \frac{C_{\parallel(\overline{\parallel})}}{t^2}\left[\ln\left(\frac{t}{\Lambda_0^2}\right)\right]^{-\gamma}, \qquad (15)$$

$$e^{\uparrow}_{\parallel(\overline{\parallel})} = \langle p_\uparrow|\mathbf{Q}_{\parallel(\overline{\parallel})}|n_\uparrow\rangle, \qquad e^{\downarrow}_{\parallel(\overline{\parallel})} = \langle p_\downarrow|\mathbf{Q}_{\parallel(\overline{\parallel})}|n_\downarrow\rangle, \qquad (16)$$

which characterize the conservation of $SU(3)$ flavor and $SU(2)$ spin symmetries in the $n \to p$ transition. Note that $\mathbf{Q}_{\parallel(\overline{\parallel})} = \sum_i Q_{\parallel(\overline{\parallel})}(i)$ with $i = 1, 2, 3$ as the chiral charge operators are coming from $Q_R \equiv J^0_R = u^\dagger_R d_R$, which convert one of the valence d quarks in $|n_{\uparrow,\downarrow}\rangle$ to be the u quark, while the converted d quark can be parallel or antiparallel to the n's helicity, denoted as the subscript ($\parallel$ or $\overline{\parallel}$). By comparing Eqs. (7) and (10) with Eqs. (13), (14), (15), and (16), we obtain

$$C_{F_1} = (e^{\uparrow}_{\parallel} + e^{\downarrow}_{\parallel})C_{\parallel} + (e^{\uparrow}_{\overline{\parallel}} + e^{\downarrow}_{\overline{\parallel}})C_{\overline{\parallel}},$$
$$C_{g_A} = (e^{\uparrow}_{\parallel} - e^{\downarrow}_{\parallel})C_{\parallel} + (e^{\uparrow}_{\overline{\parallel}} - e^{\downarrow}_{\overline{\parallel}})C_{\overline{\parallel}}, \qquad (17)$$

with $(e^{\uparrow}_{\parallel}, e^{\uparrow}_{\overline{\parallel}}, e^{\downarrow}_{\parallel}, e^{\downarrow}_{\overline{\parallel}}) = (4/3, 0, 0, -1/3)$ for the n to p transition. Similarly, we are able to relate C_{F_1} and C_{g_A} for other decay modes, given in Table I. However, C_{h_A} in Eq. (12) only has the $SU(3)$ flavor symmetry to relate different decay modes, given by

$$\langle \mathbf{B}^i_a \bar{\mathbf{B}}'^j_b|(A_\mu)^k_c|0\rangle = \bar{u}[Dd^{ijk}_{abc} + Ff^{ijk}_{abc} + Ss^{ijk}_{abc}]q_\mu\gamma_5 v, \qquad (18)$$

TABLE I. The parameters C_{F_1} and C_{g_A} in Eq. (10) are combined with $C_\parallel$ and $C_{\overline{\parallel}}$, where the upper (lower) sign is for C_{F_1} (C_{g_A}), while C_{h_A} consists of C_D, C_F, and C_S.

Matrix element	$C_{F_1}(C_{g_A})$	C_{h_A}
$\langle p\bar{p}\vert(\bar{u}u)\vert 0\rangle$	$\frac{5}{3}C_\parallel \pm \frac{1}{3}C_{\overline{\parallel}}$	$C_D + C_F + C_S$
$\langle p\bar{p}\vert(\bar{d}d)\vert 0\rangle$	$\frac{1}{3}C_\parallel \pm \frac{2}{3}C_{\overline{\parallel}}$	C_S
$\langle p\bar{p}\vert(\bar{s}s)\vert 0\rangle$	0	$C_D - C_F + C_S$
$\langle p\bar{n}\vert(\bar{u}d)\vert 0\rangle$	$\frac{4}{3}C_\parallel \mp \frac{1}{3}C_{\overline{\parallel}}$	$C_D + C_F$
$\langle \Sigma^-\bar{\Sigma}^0\vert(\bar{d}u)\vert 0\rangle$	$\frac{1}{3\sqrt{2}}(5C_\parallel \pm C_{\overline{\parallel}})$	$\sqrt{2}C_F$
$\langle \Lambda\bar{\Lambda}\vert(\bar{u}u)\vert 0\rangle$	$\frac{1}{2}C_\parallel \pm \frac{1}{2}C_{\overline{\parallel}}$	$\frac{1}{3}C_D + C_S$
$\langle \Lambda\bar{\Lambda}\vert(\bar{d}d)\vert 0\rangle$	$\frac{1}{2}C_\parallel \pm \frac{1}{2}C_{\overline{\parallel}}$	$\frac{1}{3}C_D + C_S$
$\langle \Lambda\bar{\Lambda}\vert(\bar{s}s)\vert 0\rangle$	$C_\parallel$	$\frac{4}{3}C_D + C_S$
$\langle \Lambda\bar{p}\vert(\bar{s}u)\vert 0\rangle$	$-\sqrt{\frac{3}{2}}C_\parallel$	$-\frac{1}{\sqrt{6}}(C_D + 3C_F)$
$\langle \Sigma^0\bar{p}\vert(\bar{s}u)\vert 0\rangle$	$\frac{-1}{3\sqrt{2}}(C_\parallel \pm 2C_{\overline{\parallel}})$	$\frac{1}{\sqrt{2}}(C_D - C_F)$

where $D = C_D/t^2$, $F = C_F/t^2$, and $S = C_S/t^2$ stand for the symmetric, antisymmetric, and singlet form factors for h_A; $\mathbf{B}^i_a$ and $\bar{\mathbf{B}}'^j_b$ are the baryon and antibaryon octets; and d^{ijk}_{abc}, f^{ijk}_{abc}, and s^{ijk}_{abc} are given by [27]

$$d^{ijk}_{abc} = \delta^i_b\delta^j_c\delta^k_a + \delta^i_c\delta^j_a\delta^k_b,$$
$$f^{ijk}_{abc} = \delta^i_b\delta^j_c\delta^k_a - \delta^i_c\delta^j_a\delta^k_b,$$
$$s^{ijk}_{abc} = \delta^i_b\delta^j_a\delta^k_c, \qquad (19)$$

respectively. For $\langle p\bar{n}|\bar{u}\gamma_\mu\gamma_5 d|0\rangle$, $(A_\mu)^1_2 = \bar{u}\gamma_\mu\gamma_5 d$, we obtain $C_{h_A} = C_D + C_F$ in terms of $\mathbf{B}^1_3\bar{\mathbf{B}}'^3_2 = p\bar{n}$. We also list C_{h_A} for other decay modes in Table I.

III. NUMERICAL ANALYSIS

For the numerical analysis, the CKM matrix elements and the quark masses are taken from the particle data group [9], where $m_b = 4.2$ GeV. The decay constants in Eq. (6) are given by [28,29]

$$(f_B, f_{B_s}, f_{D_s}) = (190, 225, 250)\ \text{MeV}. \qquad (20)$$

For the parameters in Table I, we refit $C_\parallel$ and $C_{\overline{\parallel}}$ by the approach of Ref. [6] with the data of $\mathcal{B}(\bar{B}^0_{(s)} \to p\bar{p})$, $\mathcal{B}(D^+_s \to p\bar{n})$, $\mathcal{B}(\bar{B}^0 \to n\bar{p}D^{*+})$, and $\mathcal{B}(\bar{B}^0 \to \Lambda\bar{p}\pi^+)$, while C_D, C_F, and C_S are newly added in the fitting. Note that the OZI suppression makes $\langle p\bar{p}|(\bar{s}s)|0\rangle = 0$, which results in $C_S = C_F - C_D$. With $N_c = 2$ fixed in a_i as the best fit, the parameters are fitted to be

$$(C_\parallel, C_{\overline{\parallel}}) = (-102.4 \pm 7.3, 210.9 \pm 85.2)\ \text{GeV}^4,$$
$$(C_D, C_F) = (-1.7 \pm 1.6, 4.2 \pm 0.7)\ \text{GeV}^4. \qquad (21)$$

As shown in Table II, we can reproduce the data of $\bar{B}^0_{(s)} \to p\bar{p}$ and $D^+_s \to p\bar{n}$. In addition, we predict the

TABLE II. The branching ratios of $B_{(s)} \to \mathbf{B}\bar{\mathbf{B}}'$ ($D_s \to \mathbf{B}\bar{\mathbf{B}}'$) decays in units of 10^{-8} (10^{-3}), where the uncertainties arise from the timelike baryonic $0 \to \mathbf{B}\bar{\mathbf{B}}'$ form factors.

Decay mode	Our result	Data
$\bar{B}^0 \to p\bar{p}$	$1.4^{+0.5}_{-0.5}$	$1.47^{+0.71}_{-0.53}$ [10]
$\bar{B}^0_s \to p\bar{p}$	$3.0^{+1.5}_{-1.2}$	$2.84^{+2.20}_{-1.69}$ [10]
$D^+_s \to p\bar{n}$	$1.3^{+13.2}_{-1.3}$	$1.30^{+0.38}_{-0.39}$ [15]
$B^- \to n\bar{p}$	$3.2^{+6.9}_{-3.0}$	⋯
$B^- \to \Lambda\bar{p}$	$3.5^{+0.7}_{-0.5}$	< 32 [30]
$\bar{B}^0 \to \Lambda\bar{\Lambda}$	$0.3^{+0.2}_{-0.2}$	< 32 [30]
$\bar{B}^0_s \to \Lambda\bar{\Lambda}$	$5.3^{+1.4}_{-1.2}$	⋯
$B^- \to \Sigma^0\bar{p}$	$5.3^{+3.8}_{-2.7}$	⋯
$B^- \to \Sigma^-\bar{\Sigma}^0$	$9.6^{+4.0}_{-3.3}$	⋯

branching ratios of $\bar{B}^0_{(s)} \to \Lambda\bar{\Lambda}$, $B^- \to (\Lambda\bar{p}, \Sigma^0\bar{p})$, and $B^- \to (n\bar{p}, \Sigma^-\bar{\Sigma}^0)$ in Table II.

IV. DISCUSSIONS AND CONCLUSIONS

When the axial-vector current is not asymptotically conserved, we can evaluate the two-body baryonic $B_{(s)}$ and D_s decays with the annihilation mechanism to explain the data. In particular, the experimental values of $\mathcal{B}(\bar{B}^0_{(s)} \to p\bar{p})$ and $\mathcal{B}(D_s^+ \to p\bar{n})$ can be reproduced. It is the violation of the PCAC that makes $\mathcal{B}(D_s^+ \to p\bar{n})$ to be of order 10^{-3}, which was considered as the consequence of the long-distance contribution in Ref. [14]. With $m_{D_s} \simeq m_p + m_{\bar{n}}$, the amplitude of $\mathcal{A}_1(D_s^+ \to p\bar{n})$ from Eq. (8) is in fact proportional to $\bar{u}(g_A + h_A)v$. Instead of $h_A = -g_A$ from the PCAC in Eq. (11) with $t = m^2_{D_s}$, our approach with $h_A = -0.7g_A$ shows that the 30% broken effect of the PCAC suffices to reveal $\mathcal{B}(D_s^+ \to p\bar{n})$. As seen from Table I, $C_{h_A} = C_D + C_F$ for the $p\bar{n}$ production with the uncertainties fitted in Eq. (21) has the solutions of $h_A = 0$ to $h_A = -g_A$, which allows $\mathcal{B}(D_s^+ \to p\bar{n}) = (0-16) \times 10^{-3}$. With the OZI suppression of $\langle p\bar{p}|(\bar{s}s)|0\rangle = 0$, which eliminates $\mathcal{A}_2$, the decay of $\bar{B}^0_s \to p\bar{p}$ is the same as that of $\bar{B}^0 \to p\bar{p}$ to be the first type. In contrast with $D_s^+ \to p\bar{n}$, since $\mathcal{A}_1(\bar{B}^0_{(s)} \to p\bar{p}) \propto m_B^2[(\frac{m_p+m_{\bar{p}}}{m_B})^2 g_A + h_A]\bar{u}\gamma_5 v$ with a suppressed g_A contribution at the m_B scale, the decay branching ratios are enhanced by h_A with m_B^2. Similarly, being of the first type, our predicted results for $\mathcal{B}(\bar{B}^0 \to \Lambda\bar{\Lambda})$, $\mathcal{B}(B^- \to n\bar{p})$, and $\mathcal{B}(B^- \to \Sigma^-\bar{\Sigma}^0)$ can be used to test the violation of the PCAC at the GeV scale.

On the contrary, $\mathcal{B}(B^- \to \Lambda(\Sigma^0)\bar{p})$ and $\mathcal{B}(\bar{B}^0_s \to \Lambda\bar{\Lambda})$ are primarily contributed from $\mathcal{A}_2$. Similar to the theoretical relation between $B^- \to p\bar{p}\ell\bar{\nu}$ [31] and $B \to p\bar{p}M$, which are associated with the same form factors in the B to $\mathbf{B}\bar{\mathbf{B}}'$ transition, resulting in the first observation of the semileptonic baryonic B decays [32], there are connections between the two-body $B^- \to \Lambda(\Sigma^0)\bar{p}$ and $\bar{B}^0_s \to \Lambda\bar{\Lambda}$ and three-body $\bar{B}^0 \to \Lambda\bar{p}\pi^+$ and $B \to \Lambda\bar{\Lambda}K$ decays with the same form factors via the (pseudo)scalar currents. As a result, without the PCAC, the observations of these two-body modes can serve as the test of the factorization, which accounts for the short-distance contribution. Note that the recent work by fitting $\bar{B}^0 \to p\bar{p}$ with the nonfactorizable contributions leads $\mathcal{B}(\bar{B}^0_s \to p\bar{p})$ and $\mathcal{B}(\bar{B}^0 \to \Lambda\bar{\Lambda})$ to be nearly zero [33], which are clearly different from our results.

In summary, we have proposed that, based on the factorization, the annihilation mechanism can be applied to all of the two-body baryonic $B_{(s)}$ and D_s decays, which indicates that the hypothesis of the PCAC is violated at the GeV scale. With the modified timelike baryonic form factors via the axial-vector currents, we are able to explain $\mathcal{B}(\bar{B}^0_{(s)} \to p\bar{p})$ and $\mathcal{B}(D_s^+ \to p\bar{n})$ of order 10^{-8} and 10^{-3}, respectively. For the decay modes that have the contributions from the (pseudo)scalar currents, they have been predicted as $\mathcal{B}(B^- \to \Lambda\bar{p}) = (3.5^{+0.7}_{-0.5}) \times 10^{-8}$, $\mathcal{B}(B^- \to \Sigma^0\bar{p}) = (5.3^{+3.8}_{-2.7}) \times 10^{-8}$, and $\mathcal{B}(\bar{B}^0_s \to \Lambda\bar{\Lambda}) = (5.3^{+1.4}_{-1.2}) \times 10^{-8}$, which can be used to test the annihilation mechanism. Besides, the branching ratios of $\bar{B}^0 \to \Lambda\bar{\Lambda}$, $B^- \to n\bar{p}$, and $B^- \to \Sigma^-\bar{\Sigma}^0$, predicted to be $(0.3, 3.2, 9.6) \times 10^{-8}$, can be viewed as the test of the PCAC, which are accessible to the experiments at the LHCb.

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