Review Article

Baryonic $B$ Meson Decays

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We review the two- and three-body baryonic $B$ decays with the dibaryon ($B\bar{B}$) as the final states. Accordingly, we summarize the experimental data of the branching fractions, angular asymmetries, and CP asymmetries. Using the $W$-boson annihilation (exchange) mechanism, the branching fractions of $B\rightarrow B\bar{B}$ are shown to be interpretable. In the approach of perturbative QCD counting rules, we study the three-body decay channels. In particular, we review the CP asymmetries of $B\rightarrow B\bar{B}M$, which are promised to be measured by the LHCb and Belle II experiments. Finally, we remark the theoretical challenges in interpreting $\mathcal{B}(B\rightarrow p\bar{p}\pi^0)$ and $\mathcal{B}(B\rightarrow p\bar{p}\nu\bar{\nu})$.

1. Introduction

The baryonic $B$ meson decays have been richly measured with the branching fractions, angular asymmetries, and CP asymmetries in two- and three-body decay channels [1–13], as summarized in Table 1. Typically, $\mathcal{B}(B\rightarrow B\bar{B})$ is as small as $10^{-8} - 10^{-7}$. Nonetheless, it is observed that $\mathcal{B}(B\rightarrow B\bar{B} M) \sim (10 - 100) \times \mathcal{B}(B\rightarrow B\bar{B})$, due to a sharply rising peak in $B\rightarrow B\bar{B} M$ observed around the threshold area of $m_{B\bar{B}} \sim m_B + m_{B\bar{B}}$ in the dibaryon invariant mass spectrum [4]. Known as the threshold effect, it enhances $\mathcal{B}(B\rightarrow B\bar{B} M)$ as large as $10^{-6}$. While the $B\bar{B}$ production shows the tendency to occur around $m_{B\bar{B}} \sim m_B + m_{B\bar{B}}$, $B\rightarrow B\bar{B}$ proceeds at $m_B$ scale, far from the threshold area. This interprets the suppressed $\mathcal{B}(B\rightarrow B\bar{B})$ [14, 15].

The partial branching fraction of $B\rightarrow B\bar{B} M$ can be a function of $\cos \theta_B$, where $\theta_B$ is the angle between the baryon and meson moving directions in the dibaryon rest frame. One hence defines the forward-backward angular asymmetry,

$$\mathcal{A}_{FB} \equiv \frac{\mathcal{B}(\cos \theta_B > 0) - \mathcal{B}(\cos \theta_B < 0)}{\mathcal{B}(\cos \theta_B > 0) + \mathcal{B}(\cos \theta_B < 0)}.$$  

In Table 1, $\mathcal{A}_{FB}(B^- \rightarrow p\bar{p}\pi^+, ppK^-) = (-40.9 \pm 3.4, 49.5 \pm 1.4)\%$ [2] indicate that one of the dibaryons favors to move collinearly with the meson.

We search for the theoretical approach to interpret the threshold effect, branching fractions, and angular asymmetries of the baryonic $B$ decays. We find that the factorization approach can be useful [16], where one factorizes (decomposes) the amplitude of the decay as two separate matrix elements. In our case, we present

$$\mathcal{M}(B\rightarrow B\bar{B}) \propto \langle B\bar{B}|\bar{q}q|0\rangle \langle 0|\bar{q}b|B\rangle,$$

$$\mathcal{M}_1(B\rightarrow B\bar{B} M) \propto \langle B\bar{B}|\bar{q}q|0\rangle \langle M|\bar{q}b|B\rangle,$$

$$\mathcal{M}_2(B\rightarrow B\bar{B} M) \propto \langle M|\bar{q}q|0\rangle \langle B\bar{B}|\bar{q}b|B\rangle.$$
Table 1: The measured branching fractions, forward-backward asymmetries ($A_{FB}$), and CP asymmetries ($A_{CP}$) for the baryonic $B$ decays, where the notation $\uparrow$ is for $A_{FB}$ with $m_{\bar{p}p} < 2.85$ GeV.

<table>
<thead>
<tr>
<th>Decay mode</th>
<th>Branching fraction</th>
<th>$A_{FB}$</th>
<th>$A_{CP}$</th>
<th>Ref.</th>
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<tbody>
<tr>
<td>$\bar{B}^0 \rightarrow \bar{p}p$</td>
<td>$(1.25 \pm 0.32) \times 10^{-8}$</td>
<td></td>
<td>$-0.41 \pm 0.11 \pm 0.03$</td>
<td>$0.04 \pm 0.07$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \Lambda \bar{\Lambda}$</td>
<td>$&lt;3.2 \times 10^{-7}$</td>
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<tr>
<td>$B^- \rightarrow \Lambda \bar{p}$</td>
<td>$(2.4^{+1.6}_{-0.9}) \times 10^{-7}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \bar{p}p$</td>
<td>$&lt;1.5 \times 10^{-8}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \bar{p}p\pi^0$</td>
<td>$(5.0 \pm 1.9) \times 10^{-7}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \bar{p}pK^0$</td>
<td>$(2.66 \pm 0.32) \times 10^{-6}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \Lambda \bar{p}\pi^+$</td>
<td>$(3.14 \pm 0.29) \times 10^{-6}$</td>
<td>$-0.41 \pm 0.11 \pm 0.03$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \Sigma^+ \bar{p}\pi^+$</td>
<td>$&lt;3.8 \times 10^{-6}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \Lambda pK^+$</td>
<td>$&lt;8.2 \times 10^{-7}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \Lambda \bar{\Lambda}K^0$</td>
<td>$(4.8^{+1.0}_{-0.9}) \times 10^{-6}$</td>
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<tr>
<td>$B^- \rightarrow \bar{p}p\pi^+$</td>
<td>$(1.62 \pm 0.20) \times 10^{-6}$</td>
<td>$(0.409 \pm 0.033 \pm 0.006)^{\uparrow}$</td>
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<tr>
<td>$B^- \rightarrow \bar{p}pK^+$</td>
<td>$(5.9 \pm 0.5) \times 10^{-6}$</td>
<td>$(0.495 \pm 0.012 \pm 0.007)^{\uparrow}$</td>
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<tr>
<td>$B^- \rightarrow \Lambda \bar{p}\pi^0$</td>
<td>$(3.0^{0.7}_{-0.7}) \times 10^{-6}$</td>
<td>$-0.16 \pm 0.18 \pm 0.03$</td>
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<tr>
<td>$B^- \rightarrow \Lambda \bar{\Lambda}\pi^+$</td>
<td>$&lt;9.4 \times 10^{-7}$</td>
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<tr>
<td>$B^- \rightarrow \Lambda \Lambda K^-$</td>
<td>$(3.4 \pm 0.6) \times 10^{-6}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \bar{p}\Lambda K^++c.c.$</td>
<td>$(5.5 \pm 1.0) \times 10^{-6}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \bar{p}pK^{*0}$</td>
<td>$(1.24^{0.28}_{-0.25}) \times 10^{-6}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \Lambda \Lambda K^{*0}$</td>
<td>$(2.5^{0.9}_{-0.9}) \times 10^{-6}$</td>
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<tr>
<td>$B^- \rightarrow \bar{p}pK^{*+}$</td>
<td>$(3.6^{0.8}_{-0.7}) \times 10^{-6}$</td>
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<tr>
<td>$B^- \rightarrow \Lambda \Lambda K^+-\pi^+\pi^-$</td>
<td>$(4.8 \pm 0.9) \times 10^{-6}$</td>
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<tr>
<td>$B^- \rightarrow \Lambda \bar{p}\phi$</td>
<td>$(8.0 \pm 2.2) \times 10^{-7}$</td>
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<tr>
<td>$B^- \rightarrow \Lambda \Lambda K^{*-}$</td>
<td>$(2.2^{+1.2}_{-0.9}) \times 10^{-6}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \Lambda \Lambda^{*-}$</td>
<td>$(1.54 \pm 0.18) \times 10^{-5}$</td>
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<tr>
<td>$B^- \rightarrow \Sigma_+^0\bar{p}$</td>
<td>$&lt;2.4 \times 10^{-5}$</td>
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<tr>
<td>$B^- \rightarrow \Sigma_0^0\bar{p}$</td>
<td>$(2.9 \pm 0.7) \times 10^{-5}$</td>
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<tr>
<td>$B^- \rightarrow \bar{p}pD^+$</td>
<td>$&lt;1.5 \times 10^{-5}$</td>
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<tr>
<td>$B^- \rightarrow \bar{p}pD^{*-}$</td>
<td>$&lt;1.5 \times 10^{-5}$</td>
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<tr>
<td>$B^- \rightarrow \Lambda \bar{p}D^0$</td>
<td>$(1.43 \pm 0.32) \times 10^{-5}$</td>
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<tr>
<td>$B^- \rightarrow \Lambda \bar{p}D^{*0}$</td>
<td>$&lt;5 \times 10^{-5}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow n\bar{p}D^{**}$</td>
<td>$(1.4 \pm 0.4) \times 10^{-3}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \bar{p}pD^0$</td>
<td>$(1.04 \pm 0.07) \times 10^{-4}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \bar{p}pD^{*0}$</td>
<td>$(0.99 \pm 0.11) \times 10^{-4}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \Lambda \bar{p}D^*_0$</td>
<td>$(2.8 \pm 0.9) \times 10^{-5}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \Lambda \bar{p}D^*$</td>
<td>$(2.5 \pm 0.4) \times 10^{-5}$</td>
<td>$-0.08 \pm 0.10$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \Lambda \bar{p}D^{**}$</td>
<td>$(3.4 \pm 0.8) \times 10^{-5}$</td>
<td>$+0.55 \pm 0.17$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \Lambda \Lambda \bar{D}^0$</td>
<td>$(1.00^{+0.30}_{-0.26}) \times 10^{-5}$</td>
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<tr>
<td>$\bar{B}^0 \rightarrow \Sigma^0 \Lambda \bar{D}^0+c.c.$</td>
<td>$&lt;3.1 \times 10^{-5}$</td>
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</table>
where \((q\bar{q}')\) and \((q\bar{b})\) stand for the quark currents and the matrix element of \(\langle BB' \mid \bar{q}q' \rangle \langle 0 \mid (BB' \mid \bar{q}b)B \rangle\) can be parameterized as the timelike baryonic \((B \to BB'\) transition) form factors \(F_{BB'}\). Moreover, one derives \(F_{BB'} \propto 1/n^2\) in perturbative QCD (pQCD) counting rules [17–24], where \(f = (p_B + p_B')^2\) and \(n\) accounts for the number of the gluon propagators that attach to the baryon pair. It results in \(d\sigma/dm_{BB'} \propto 1/n^2\), which shapes a peak around \(m_{BB'} \sim m_B + m_{B'}\) in the \(m_{BB'}\) spectrum, and then, the threshold effect can be interpreted. In the \(B \to pp\) transition, there exists the term of \((p_B - p_p) \bar{u}(y_z)\nu\) for \(F_{BB'}\) [24], which is reduced as \(E_p - E_B \bar{u}(y_z)\nu\) in the \(pp\) rest frame. Since \(E_p - E_B \propto \cos \theta_p\), the term for \(F_{BB'}\) can be used to describe the highly asymmetric \(\mathcal{A}_{FB}(B \to pp\pi^+\pi^-)\). Alternatively, the baryonic B decays is studied with the pole model, where the nonfactorizable contributions can be taken into account [25–28].

We have explained \(\mathcal{A}(B \to BB') \sim 10^{-8} - 10^{-7}\) [29]. We have studied \(B \to BB'M\) and explained the branching fractions and CP asymmetries [30–38]. In addition, we have predicted \(\mathcal{A}(B^0 \to p\bar{A}K^- + \Lambda\bar{p}K^+) = (5.1 \pm 1.1) \times 10^{-6}\) [37], in excellent agreement with the value of \((5.46 \pm 0.61 \pm 0.57 \pm 0.50 \pm 0.32) \times 10^{-6}\) measured by LHCb [8]. This demonstrates that the theoretical approach can be reliable. Therefore, we would like to present a review, in order to illustrate how the approach of pQCD counting rules based on the factorization can be applied to the baryonic B decays. We will also review our theoretical results that have explained the branching fractions of \(B \to BB'\) and \(B \to BB'M\), particularly, the CP asymmetries, promising to be observed by future measurements.

2. Formalism

To review the two-body baryonic B decays, we take \(B^0 \to pp\) as our example. According to Figure 1, \(B^0 \to pp\) is regarded as a \(W\)-boson exchange process [29, 39–41]. In the factorization, we derive the amplitude as [29]

\[
\mathcal{A}(B^0 \to pp) = \frac{G_F}{\sqrt{2}} V_{ub} V_{ut} a_1 \langle pp \mid \bar{u}p(1 - y_s)u \mid 0 \rangle \langle 0 \mid \bar{d}q(1 - y_s)b \mid \bar{p}p \rangle,
\]

where \(G_F\) is the Fermi constant and \(V_{ub}^{(*)}\) the Cabibbo-Kobayashi-Maskawa (CKM) matrix element. One has defined \(\langle 0 \mid \bar{d}q(1 - y_s)b \mid B^0 \rangle = -(iF_{B^0}q_{\mu})\) for the \(B^0\) meson annihilation, where \(f_{B^0}\) is the decay constant and \(q_{\mu}\) the four-momentum. For the \(pp\) production, the matrix elements read [22, 23]

\[
\langle BB' \mid V_{\mu} \rangle \approx \bar{u}(F_1 y_{\mu} + \frac{F_2}{m_B + m_{B'}} + i\alpha_{\mu} q_{\mu} A_{\mu}) \nu,
\]

\[
\langle BB' \mid A_{\mu} \rangle \approx \bar{u}(g_A y_{\mu} + \frac{h_A}{m_B + m_{B'}} q_{\mu}) \gamma_5 \gamma_{\mu},
\]

with the (axial-)vector current \(V(A)_{\mu} = \bar{q}g_{\mu}^{(A)}(y_s)q'\), where \(F_{1,2}, g_A\), and \(h_A\) are the timelike baryonic form factors.

At a very large momentum transfer \((Q^2 \to \infty)\), the approach of pQCD counting rules results in [17–21]

\[
F_{1,2}, g_A \propto \frac{\alpha_s^2(Q^2)}{Q^4} \ln \left(\frac{Q^2}{\Lambda^2}\right)^{-4(3\beta)},
\]

where \(\beta = 11 - 2n_f/3\) is the \(\beta\) function of QCD to one loop, \(n_f = 3\) the flavor number, and \(\Lambda = 0.3\) GeV the scale factor. Moreover, \(\alpha_s(Q^2) \approx (4\pi/\beta)[\ln(Q^2/\Lambda^2)]^{-1}\) is the running coupling constant in the strong interaction [20]. Interestingly, \(\alpha_s^2/Q^4\) reflects the fact that one needs two hard gluon propagators to attach to the baryons as drawn in Figure 1 [42], whereas \(\ln(Q^2/\Lambda^2)^{-4(3\beta)}\) is caused by the wave function.

As \(V_{\mu}\) and \(A_{\mu}\) are combined as the right- or left-handed chiral current, that is, \(J_{\mu}^{B\ell} = (V_{\mu} \pm A_{\mu})/2\), one obtains \(\langle B \mid J_{\mu}^{\ell R}(\ell_{B}) \mid B_{\ell R} \rangle\) for the space-like \(B \to B'\) transition. With the right-handed current, the matrix elements can be written as [19, 29]

\[
\langle B_{\ell R} \mid J_{\mu}^{\ell R}(\ell_{B}) \mid B \rangle = \bar{u}(g_{\mu} - \frac{1 + y_{\mu}}{2} F_{R} + \frac{1 - y_{\mu}}{2} F_{L}) u,
\]

where \(\mathcal{Q}_{\mu}^{\ell R}(\ell_{B}) = \langle B_{\ell R} \mid J_{\mu}^{\ell R}(\ell_{B}) \rangle\) and \(F_{R,L}\) are the chiral form factors. With \(q_{\ell}(i = 1, 2, 3)\) denoting one of the valence quarks in \(B\), \(Q = \sum_{\mu=0}^{3} \) known as the chiral charge is able to change the flavor for \(q_{\ell}\), such that \(B\) is transformed as \(B'\). Note that the chirality is regarded as the helicity at \(Q^2 \to \infty\). Since the helicity of \(q_{\ell}\) can be (anti-)parallel \((\|\|)\) to the helicity of \(B\), we define \(Q_{\|}^{\ell R}(i)\) that is responsible for acting on \(q_{\ell}\). Thus, the approach of pQCD counting rules leads to [19, 29]

\[
F_R = e_R^0 F_R + e_R^1 F_R,\quad F_L = e_L^1 F_L + e_L^2 F_L,
\]

with \(e_R^0 = \langle B_{\ell R} \mid Q_{\|}^{\ell R}(i) \mid B_{\ell R} \rangle,\)

\(e_R^1 = \langle B_{\ell R} \mid Q_{\|}^{\ell R}(i) \mid B_{\ell R} \rangle,\) and \(Q_{\|}^{\ell R}(i) = \sum_{\ell} Q_{\|}^{\ell R}(i)\), where the \(SU(3)\) flavor \((SU(3)_f)\) and \(SU(2)\) spin symmetries are both respected.

![Figure 1: Feynman diagram for $B^0 \to pp$.](image-url)
symmetry, the spacelike form factors behave as the timelike ones, such that one can relate $F_1$ and $g_A$ with the chiral form factors in Equations (6) and (7) derived in the spacelike region, leading to $F_1(g_A) = (e_0^B + e_1^B) F_1 + (e_0^B - e_1^B) F_2$. In addition to the momentum dependence of Equation (5), $F_1$ and $g_A$ are presented as [22, 23, 29]

$$ F_1 = \frac{C_{F_1}}{t^2} \ln \left( \frac{t}{\Lambda^2} \right)^\gamma, \quad g_A = \frac{C_{g_A}}{t^2} \ln \left( \frac{t}{\Lambda^2} \right)^\gamma, $$

where $\gamma = 2 + 4/(3\beta) = 2.148$.

For $\langle pp|\bar{u}u|0\rangle$, we obtain [23]

$$ C_{F_1} = \frac{5}{3} C_1 + \frac{1}{3} C_2, $$

$$ C_{g_A} = \frac{5}{3} C_1 - \frac{1}{3} C_2, $$

where $C_{[1]}$ is from $F_{[1]} \equiv C_{[1]} / t^2 \ln \left( t/\Lambda^2 \right)^\gamma$, and we have used $(e_0^B, e_1^B) = (5/3, 0)$ and $(e_0^B, e_2^B) = (0, 1/3)$ [23, 29]. In Ref. [43], the pQCD calculation causes $F_2 = F_1 / (t \ln [t/\Lambda^2])$, indicating that $F_2$ has a suppressed contribution. Taking the form factors as the inputs, we reduce the amplitude of $\bar{B}^0 \rightarrow \bar{p}\bar{p}$ as

$$ M \propto \frac{1}{(m_p + m_p)} \bar{u} \left[ (m_p + m_p) g_A + m_p^2 h_A \right] \gamma_{\bar{p}V}, $$

where $F_{1(2)}$ has been vanishing, in accordance with the conservation of vector current (CVC) $q^\mu \langle pp|\bar{u}\gamma_\mu u|0\rangle = 0$.

Using the partial conservation of axial-vector current (PCAC), where $q^\mu (\bar{B}B|^A_\mu|0) = 0$, it is obtained that [23, 29, 39–41, 44]

$$ h_A = \frac{(m_b + m_B)^2 g_A}{t - m_M^2}, $$

by which $g_A$ and $h_A$ cancel each other and then $\mathcal{M} = 0$. This seems that the $W$-exchange (annihilation) mechanism based on the factorization fails to explain $\mathcal{B}(B \rightarrow \bar{B}\bar{B}^\rightarrow)$. As a consequence, one turns to think of the nonfactorizable effects as the main contributions [25, 44–50].

In Equation (11), $\mathcal{M}(B \rightarrow \bar{B}\bar{B}^\rightarrow)$ describes a meson pole, so that $B \rightarrow \bar{B}\bar{B}^\rightarrow$ can be regarded to receive the contribution from the intermediate process of $B \rightarrow M \rightarrow \bar{B}\bar{B}^\rightarrow$, which is much suppressed. On the other hand, there might exist a QCD-based contribution to $h_A$, by which $\mathcal{M}(B \rightarrow \bar{B}\bar{B}^\rightarrow) \neq 0$, and PCAC is violated. Here, we choose to parameterize $h_A$ with slightly violated PCAC in the timelike region. To this end, we derive $h_A + g_A = 0$ with $q^\mu (\bar{B}B|^A_\mu|0) = 0$ at the threshold area of $t = (m_p + m_p)^2$, where the meson pole is supposed to be inapplicable [40, 41]. Since the QCD-based calculation of $h_A$ is still lacking, besides $h_A + g_A = 0$ suggesting $h_A \propto g_A^2$, we are allowed to present $h_A = C_{h_A}/t^2$ for its momentum dependence [29].

To describe the three-body baryonic $B$ decays, we take $B \rightarrow \bar{p}\bar{p}V$ with $V = p^{-(0)}$ or $K^{-(0)}(K^{+})$ as our examples. According to Figure 2, the amplitudes are given by [30–32]

$$ \mathcal{M}(B \rightarrow \bar{p}\bar{p}V) = \frac{C_E}{\sqrt{2}} \alpha_V \langle V | q_{\gamma_\mu} (1 - \gamma_5) q | 0 \rangle \langle \bar{p}\bar{p} | \gamma_\mu (1 - \gamma_5) b | B \rangle. $$
In Equation (12), one presents that $$V_\mu^0 q^0 \langle 0 | V_\mu^0 q^0 | 0 \rangle$$ = $f_V m_V e^*_\mu$, where $f_V$ and $e^*_\mu$ are, respectively, the decay constant and polarization four-vector of the vector meson. The amplitudes are both associated with the matrix elements of the $B \to B^{\pm} \to B^{0} \to B^{0} \to B^{0} \to B^{0} \to B^{0} \to B^{0}$ transition, and we parameterize them as [24]
where \( V_{\mu}^8(A^h_{\mu}) = i g_\mu^3(y_g^8)b, \) \( p = p_B - (p_B + p_B), \) and \( \langle f, g_i \rangle \) \((i = 1, 2, \ldots, 5)\) are the B \( \rightarrow \bar{B}B \) transition form factors. Inspired by pQCD counting rules [17, 18, 21, 23, 24], the momentum dependences for \( f_i \) and \( g_i \) are given by

\[
 f_i = \frac{D_{i}^f}{m^3}, \quad g_i = \frac{D_{i}^g}{m^3},
\]

with \( D_{i}^f, D_{i}^g \) to be extracted by the data. According to the gluon lines in Figure 2, \( n \) in \( 1/n^2 \) should be \( 2 + 1 \), which accounts for two gluon propagators attaching to the valence quarks in \( \bar{B}B \) and an additional one for kicking (speeding up) the spectator quark in B [23]. For the gluon kicking, it is similar to the meson transition form factor derived as \( F_M \propto 1/q^2 \) in pQCD counting rules [20, 51], where \( 1/q^2 \) is for a hard gluon to transfer the momentum to the spectator quark in the meson.

Like the case of \( F_i \) and \( g_A \), we relate \( \langle f_i, g_i \rangle \) to the B \( \rightarrow \bar{B}B \) chiral form factors, which is in terms of [23, 24]

\[
\langle B_{\mu \nu}(\mu) \rangle = \frac{1}{2} \left[ 1 + \frac{y_g^8}{2} G_{G} + \frac{1 - y_g^8}{2} G_{L} \right] u_{i} + i \bar{y}_{\mu} p_{\mu} \left[ 1 + \frac{y_g^8}{2} G_{G} + \frac{1 - y_g^8}{2} G_{L} \right] u_{i},
\]

where \( |B_{\mu} \rangle \) \( \rightarrow |\bar{b}g \rangle \) has been used. In addition, we obtain

\[
 G_{R(L)} = e_{R(L)} G_{G} + e_{R(L)} G_{L} \quad \text{and} \quad G_{R(L)} = e_{R(L)} G_{G} + e_{R(L)} G_{L},
\]

which are similar to \( F_{R(L)} \) in Equation (7). Under the SU(3)

\[
\text{flavor and SU}(2) \text{ spin symmetries, together with } G_{R(L)} = D_{R(L)} \text{, it is derived that [23, 24]}
\]

\[
\begin{align*}
 D_{g_i} &= \frac{5}{3} D_{i} - \frac{1}{3} D_{j}, \\
 D_{g_j} &= \frac{5}{3} D_{j} - \frac{1}{3} D_{i}, \\
 D_{g_i} &= \frac{1}{3} D_{i} - \frac{2}{3} D_{j}, \\
 D_{g_j} &= \frac{1}{3} D_{j} - \frac{2}{3} D_{i}, \\
 D_{g_i} &= \frac{-1}{3} D_{i} - \frac{2}{3} D_{j},
\end{align*}
\]

for \( \langle p\bar{p}|\bar{b}g|B^- \rangle \) and \( \langle p\bar{p}|\bar{b}g|B^0 \rangle \), respectively. We also review the direct CP asymmetry, defined by

\[
\mathcal{A}_{CP}(B \rightarrow \bar{B}B') = \frac{\Gamma(B \rightarrow \bar{B}B') - \Gamma(B \rightarrow B'B')} {\Gamma(B \rightarrow \bar{B}B') + \Gamma(B \rightarrow B'B')},
\]

where \( \Gamma \) denotes the decay width and \( B \rightarrow \bar{B}B' \) the antiparticle decay.
3. Numerical Analysis

For the numerical analysis, we adopt the CKM matrix elements as [1]

\[ V_{ub} = A \lambda^3 (\rho - i \eta), \quad V_{ud} = 1 - \lambda^2/2, \quad V_{us} = \lambda, \]

\[ V_{tb} = 1, \quad V_{td} = A \lambda^2, \quad V_{ts} = -A \lambda^2, \]

with \( \lambda = 0.22453 \pm 0.00044, \quad A = 0.836 \pm 0.015, \quad \rho = 0.12 \pm 0.018, \) and \( \eta = 0.355 \pm 0.011 \) in the Wolfenstein parameterization, where \((\rho, \eta) = (1 - \lambda^2/2) \times (\rho, \eta).\) The decay constants are given by \((f_B, f_{B^*}) = (0.19, 0.21, 0.22) \) GeV [1, 52].

With the global fit to the data, we obtain [29, 37]

\[
\begin{align*}
(C_1, C_2 | C_{h_1}) &= (-102.4 \pm 7.3, 210.9 \pm 85.2, 8.4 \pm 1.4) \text{ GeV}^4, \\
(D_1, D_2) &= (45.7 \pm 33.8, -298.2 \pm 34.0) \text{ GeV}^5, \\
(D_1^2, D_1^3, D_2^3) &= (33.1 \pm 30.7, -203.6 \pm 133.4, 6.5 \pm 18.1, -147.1 \pm 29.3) \text{ GeV}^4.
\end{align*}
\]

We take \( a_i \equiv c_i^{\text{eff}} + c_i^{\text{eff}} / N_c \) for \( i = \text{odd} \) (even) with \( N_c \) the color number, where the effective Wilson coefficients \( c_i \) come from Ref. [16]. For \( a_i \) in \( B^0 \rightarrow p\bar{p} \rho^0 \) and \( B^0 \rightarrow p\bar{p} \), we use \( N_c = 2 \) to take into account the nonfactorizable QCD corrections. We can hence present our theoretical calculations for \( B^0 \rightarrow p\bar{p} (V) \) in Table 2, along with the other \( B^0 \rightarrow B B' M \) results. Besides, we present our studies of the angular and CP asymmetries.

4. Discussions and Conclusions

The theoretical results in Table 2 can agree well with the data, which is based on the factorization, pQCD counting rules, and baryonic form factors. In particular, several CP asymmetries are predicted as large as 10-20%, promising to be measured by LHCb and Belle II [55, 56]. It is reasonable to extend the theoretical approach to \( B^0 \rightarrow B B' M \) and \( B^0 \rightarrow B B' M \) [34, 38, 57], where \( B_i (M_i) \) denotes a baryon (meson) containing a charm quark. As a consequence, \( \mathcal{B}(B^0 \rightarrow p\bar{p} \eta^0) \) can also be explained (see Table 3). Nonetheless, \( \mathcal{B}(B^0 \rightarrow p\bar{p} \rho^-) = (28.8 \pm 2.1) \times 10^{-6} \) and \( \mathcal{B}(B^0 \rightarrow p\bar{p} \eta^-) = (1.04 \pm 0.24 \pm 0.12) \times 10^{-7} \) have been predicted not verified by the observations [1, 32, 58, 59],

\[
\begin{align*}
\mathcal{B}(B^0 \rightarrow p\bar{p} \rho^- \pi^0) &= (4.6 \pm 1.3) \times 10^{-6}, \\
\mathcal{B}(B^0 \rightarrow p\bar{p} \eta^- \bar{\nu}_\mu) &= (5.27^{+0.23}_{-0.24} \pm 0.21 \pm 0.15) \times 10^{-6},
\end{align*}
\]

where the amplitude of \( B^0 \rightarrow p\bar{p} \eta^0 \bar{\nu}_\mu \) is given by

\[
\mathcal{M}(B^0 \rightarrow p\bar{p} \eta^0 \bar{\nu}_\mu) = \frac{G_F}{\sqrt{2}} V_{ub} \Big( \frac{p}{b} | i \nu_\mu (1 - \gamma_5) b | B^0 \Big) \bar{\nu}_\mu (1 - \gamma_5) \nu_\mu.
\]


[54] Y. K. Hsiao, Angular asymmetry in charmless $B \to \bar{p} p M$ decays in preparation.


