Direct CP violation in internal $W$-emission dominated baryonic $B$ decays

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Abstract The observation of CP violation has been experimentally verified in numerous $B$ decays but is yet to be confirmed in final states with half-spin particles. We focus our attention on baryonic $B$-meson decays mediated dominantly through internal $W$-emission processes and show that they are promising processes to observe for the first time the CP violating effects in $B$ decays to final states with half-spin particles. Specifically, we study the $\bar{B}^0 \to p \bar{p} \pi^0(\rho^0)$ and $\bar{B}^0 \to p \bar{p} \pi^+\pi^−$ decays. We obtain $B(\bar{B}^0 \to p \bar{p} \pi^0) = (5.0 \pm 2.1) \times 10^{-7}$, in agreement with current data, and $B(\bar{B}^0 \to p \bar{p} \rho^0) \simeq B(\bar{B}^0 \to p \bar{p} \pi^0)/3$. Furthermore, we find $A_{CP}(\bar{B}^0 \to p \bar{p} \pi^0, p \bar{p} \rho^0, p \bar{p} \pi^+\pi^−) = (−16.8 ± 5.4, −12.6 ± 3.0, −11.4 ± 1.9)\%$. With measured branching fractions $B(\bar{B}^0 \to p \bar{p} \pi^0, p \bar{p} \pi^+\pi^−) \sim O(10^{-6})$, we point out that $A_{CP} \sim −(10 − 20)\%$ can be new observables for CP violation, accessible to the Belle II and/or LHCb experiments.

1 Introduction

The investigation of CP violation (CPV) has been one of the most important tasks in hadron weak decays. In the Standard Model (SM), CPV arises from a unique phase in the Cabibbo-Kobayashi-Maskawa (CKM) quark mixing matrix; however, it is insufficient to explain the matter and antimatter asymmetry of the Universe. To try and shed light on solving the above puzzle, a diverse set of observations related to CPV is necessary. So far, direct CP violation has been observed in $B$ and $D$ decays [1,2]. With $Re(\epsilon'/\epsilon)$, it is also found in kaon decays [3]. Although the decays involving half-spin particles offer an alternative route, evidence for CP violation is not richly provided [4,5].

Baryonic $B$ decays can be an important stage to investigate CPV within the SM and beyond. With $M^{(s)}$ denoting a pseudoscalar (vector) meson such as $K^{(s)}$, $\pi$, $\rho$, $D^{(s)}$, the $B \to p \bar{p} M^{(s)}$ decays have been carefully studied by the $B$ factories and the LHCb experiment [5–11]. Experimental information includes measurements of branching fractions, angular distribution asymmetries, polarization of vector mesons in $B \to p \bar{p} K^*$, Dalitz plot information, and $p \bar{p}$ ($M^{(s)} p$) invariant mass spectra. This helps to improve the theoretical understanding of the di-baryon production in $B \to BB'M$ [12–16], such that the data can be well interpreted. Predictions are confirmed by recent measurements. For example, one obtains $B(\bar{B}^0 \to p\Lambda K^− + \Lambda \bar{p} K^+) = (5.1 ± 1.1) \times 10^{-6}$ [17], in excellent agreement with the value of $(5.46 ± 0.61 ± 0.57 ± 0.32 ± 0.32) \times 10^{-6}$ measured by LHCb [18]. Moreover, the theoretical extension to four-body decays allows to interpret $B(\bar{B}^0 \to p \bar{p} \pi^−\pi^−)$ [19–21]. The same can be said for CP asymmetries.

In this report we focus our attention on the baryonic $B$-meson decays mediated dominantly through the internal $W$-emission diagrams. Although the internal $W$-emission decays are regarded as suppressed processes, the measured branching fractions of the baryonic $B$ decays

\begin{align}
B(\bar{B}^0 \to p \bar{p} \pi^0) &= (5.0 \pm 1.8 \pm 0.6) \times 10^{-7}, \\
B(\bar{B}^0 \to p \bar{p} \pi^+\pi^-) &= (2.7 \pm 0.1 \pm 0.1 \pm 0.2) \times 10^{-6},
\end{align}

are not small [19,22], which make these modes an ideal place to observe for the first time CP violation in $B$ decays to final states with half-spin particles. Therefore, we will study the branching fractions for the decays of $\bar{B}^0 \to p \bar{p} \pi^0(\rho^0)$, $p \bar{p} \pi^+\pi^−$, and predict their direct CP violating asymmetries.
The effective Hamiltonian is given by \[ H \]

\[ V_{ij} \]

where \( G_F \) is the Fermi constant, \( c_{i \langle j} \) are the Wilson coefficients, and \( V_{ij} \) are the CKM matrix elements. The four-quark operators \( O_{i \langle j} \) for the tree (penguin)-level contributions are written as

\[ O_{1} = (\tilde{d}_a u_\beta)_{V-A} (\bar{u}_\beta b_\alpha)_{V-A} , \]

\[ O_{2} = (\tilde{d}_a u_\beta)_{V-A} (\bar{u}_\beta b_\alpha)_{V-A} , \]

\[ O_{3\langle5} = (\tilde{d}_a b_\alpha)_{V-A} \sum_q (\tilde{q}_\beta q_\beta)_{V=\pm A} , \]

\[ O_{4\langle6} = (\tilde{d}_a b_\beta)_{V-A} \sum_{\bar{q}_\beta q_\beta} (\bar{q}_\bar{q} q_\bar{q})_{V=\pm A} , \]

\[ O_{7\langle9} = \frac{3}{2} (\tilde{d}_a b_\alpha)_{V-A} \sum_q e_q (\bar{q}_\beta q_\beta)_{V=\pm A} , \]

\[ O_{8\langle10} = \frac{3}{2} (\tilde{d}_a b_\beta)_{V-A} \sum_q e_q (\bar{q}_\bar{q} q_\bar{q})_{V=\pm A} . \]

where \( q = (u, d, s) \), \( (\bar{q}_1 q_2)_{V=\pm A} = \bar{q}_1 \gamma_\mu (1 \pm \gamma_5) q_2 \), and the subscripts \((\alpha, \beta)\) denote the color indices. With the identity of \( \delta_{\beta\beta'} \delta_{\alpha\alpha'} = \delta_{\alpha\beta} \delta_{\alpha'\beta'} / N_c + 2 T^{a}_{\alpha} T^{a}_{\beta}, \) where \( N_c = 3 \) is the color number, \( O_1 \) and \( O_{1+1} \) can be related. For example, we have \( O_1 = O_2 / N_c + 2 \bar{d} \gamma_\mu (1 - \gamma_5) T^a \bar{u} u_\mu (1 - \gamma_5) T^a b \) with \( T^a \) the Gell-Mann matrices.

In the factorization ansatz \([24,25]\), one is able to express \( \langle h_1 h_2 | O \rangle B \) as a product of two factors, \( \langle h_1 | J_1 \rangle 0 \) and \( \langle h_2 | J_2 | B \rangle \), where \( O = J_1 \cdot J_2 \) is the product of the two color singlet quark currents \( J_1 \) and \( J_2 \) and \( h_{1,2} \) denote the hadron states. The matrix elements \( \langle h_1 | J_1 \rangle 0 \) and \( \langle h_2 | J_2 | B \rangle \) are obtained in such a way that the flavor quantum numbers of \( J_{1,2} \) match the hadron states in the separate matrix elements. We hence decompose \( \langle p \bar{p} \pi^0 | O_2 | B \rangle \) as \([15,16]\)

\[ \langle O_2 \rangle_a = \langle \pi^0 | (\bar{u}_\beta u_\alpha)_{V-A} | 0 \rangle \langle p \bar{p} | (\bar{d}_a b_\alpha)_{V-A} | B \rangle , \]

\[ \langle O_2 \rangle_d = \langle p \bar{p} | (\bar{u}_\beta u_\alpha)_{V-A} | 0 \rangle \langle \pi^0 | (\bar{d}_a b_\alpha)_{V-A} | B \rangle . \]

where the Fierz reordering has been used to exchange \( \bar{d}_a, \bar{u}_\beta \). The amplitudes \( \langle O_2 \rangle_{a,d} \) correspond to the two configurations depicted in Fig. 1a, d, respectively. As depicted in Fig. 2 for the \( b \rightarrow u \bar{d} \) transition, dynamically, the \( d \)-quark moves collinearly with the spectator quark \( \bar{d} \) from \( B_0(bd) \), so that in Fig. 1d the \( dd \) for the \( p \bar{p} \) formation can...
be seen as a consequence of the $B$ meson transition, which is in accordance with the matrix element of $\langle p\bar{p}((db)\bar{B})^0\rangle$. Moreover, since $u\bar{u}$ and $d\bar{d}$ in the $B^0$ rest frame can be seen to move in opposite directions, we take $\pi^0(u\bar{u})$ in Fig. 1d as the recoiled state, in accordance with $\langle \pi^0| (u\bar{u})|0\rangle$ representing the vacuum. On the other hand, $\langle p\bar{p}\pi^0|O_1|B^0\rangle$ is expressed as $\langle O_1|_{u(d)} = \langle O_2|_{u(d)} / N_c + \langle \chi_1 \rangle$ with $\langle \chi_1 \rangle \equiv \langle p\bar{p}\pi^0|2\pi Y_1\rangle(1 - \gamma_5)T^a\pi Y_1(1 - \gamma_5)T^a|B^0\rangle$. The $T^a$ in $\langle \chi_1 \rangle$ correspond to the gluon exchange between the two currents, which causes an inseparable connection between the final states. Hence, $\langle \chi_1 \rangle$ is regarded as the non-factorizable QCD corrections. Subsequently, we note that $\langle p\bar{p}\pi^0|c_1O_1 + c_2O_2|B^0\rangle = a_2\langle O_2|_{u,d} + a_2 = c_2^{\text{eff}} + c_2^{\text{eff}} / N_c$, where $c_2^{\text{eff}}$ represents the effective Wilson coefficient for $c_1$ to receive the next-to-leading-order contributions [25]. In the generalized edition of the factorization, one varies $N_c$ between 2 and infinity in order to estimate $\langle \chi_1 \rangle$ [15, 24, 25]. This makes $N_c$ a phenomenological parameter determined by data.

To complete the amplitudes, we extend our calculation for $\langle p\bar{p}\pi^0|c_1O_1 + c_2O_2|B^0\rangle$ to the penguin-level diagrams, as depicted in Fig. 1b, c, e, f. Moreover, with $\pi^0$ replaced by $\rho^0$ and $\pi^+\pi^-$, we get the amplitudes of $B^0 \to p\bar{p}\rho^0$ and $\bar{B}^0 \to p\bar{p}\pi^+\pi^-$, respectively. Hence, the decay amplitudes of $\bar{B}^0 \to p\bar{p}X_M$ with $X_M \equiv (\pi^0(\rho^0), \pi^+\pi^-)$ can be written as [16, 17, 20]

$$A(\bar{B}^0 \to p\bar{p}X_M) = A_1(X_M) + A_2(X_M),$$

with $A_{1,2}(X_M)$ corresponding to Fig. 1a–c and d–f, respectively. Explicitly, $A_{1,2}$ are given by [15–17, 25–27]

$$A_1(X_M) = \frac{G_F}{\sqrt{2}} \left[ (p\bar{p}|d\gamma^\mu(\alpha_2^+ - \alpha_2^- \gamma_5)d|0) \right.\right.$$

$$\left. + (\langle p\bar{p}|d\gamma^\mu(\alpha_2^- - \alpha_2^+ \gamma_5)d|0\rangle \right) \times \langle X_M|d\gamma_1(1 - \gamma_5)b|\bar{B}^0\rangle$$

$$\left. + a_6\langle p\bar{p}|d\gamma_1(1 + \gamma_5)b|\bar{B}^0\rangle \right) \times \langle X_M|d\gamma_1(1 - \gamma_5)b|\bar{B}^0\rangle \right),$$

$$A_2(X_M) = \frac{G_F}{\sqrt{2}} \left[ \langle X_M|d\gamma^\mu(\alpha_2^+ - \alpha_2^- \gamma_5)d|0\rangle \right.$$

$$\left. + (\langle p\bar{p}|d\gamma^\mu(\alpha_2^- - \alpha_2^+ \gamma_5)d|0\rangle \right) \times \langle X_M|d\gamma_1(1 - \gamma_5)b|\bar{B}^0\rangle$$

$$\left. + a_6\langle X_M|d\gamma_1(1 + \gamma_5)b|\bar{B}^0\rangle \right) \times \langle p\bar{p}|d\gamma_1(1 - \gamma_5)b|\bar{B}^0\rangle \right).$$

The parameters $a_i$ are defined as

$$a_2^\pm = V_{ub}V_{ub}^*a_2 - V_{td}V_{td}^*(a_3 \pm a_5 \pm a_7 + a_9),$$

$$a_3^\pm = -V_{ub}V_{ub}^* \left( a_3 + a_5 \pm a_7 \pm \frac{a_9}{2} - \frac{a_{10}}{2} \right),$$

$$a_6 = V_{ub}V_{ub}^*(2a_6 - a_8),$$

with $a_i \equiv c_i^{\text{eff}} + c_i^{\text{eff}} / N_c$ for $i \text{ even}$ [25]. We note that $A_2(\pi^+\pi^-)$ is neglected since $A_1(\pi^+\pi^-) \gg A_2(\pi^+\pi^-)$ [20].

The $B \to X_M$ transition matrix elements in $A_1(X_M)$ are written as [28–31]

$$\langle M(p)|q\gamma_\mu b|B(p_B)\rangle = \frac{1}{2V_1} \left[ (p_B + p)^\mu - \frac{m_B^2 - m_M^2}{q^2} q^\mu \right] F_{BM}^{BM}$$

$$+ \frac{m_B^2 - m_M^2}{q^2} q^\mu F_{BM}^{BM},$$

$$\langle M^*(p)|q\gamma_\mu s b|B(p_B)\rangle = \epsilon_{\mu\nu\rho\sigma} q^\mu p_B^\nu p^\rho B^\sigma 2V_1 \left( m_B + m_M^* \right),$$

$$\langle M^*(p)|q\gamma_\mu s b|B(p_B)\rangle$$

$$= i \left[ \epsilon_{\mu\nu\rho\sigma} q^\mu (m_B^2 - m_M^{*2}) A_1$$

$$\left. + i \frac{q^\mu (2m_M^*) A_0 \right) \right.$$}

$$\langle M_1(p_1)M_2(p_2)\rangle = \epsilon_{\mu\nu\rho\sigma} q^\mu (2p_2 - p_1)^\rho (p_2 - p_1)^\sigma h$$

$$+ i w_+ (p_2 + p_1) + i w_- (p_2 - p_1) + ir q^\mu, \quad (8)$$

where $\epsilon_{\mu\nu\rho\sigma}$ is the polarization vector of $M^*, q_{\mu} = (p_B - p)^{\mu} = (p_B - p - p_1 - p_2)^{\mu}$ as the momentum transfer for the $B \to X_M$ transition, $(F_{BM}^{BM}, V_1, A_{0,1,2})$ the $B \to M^*(*)$ transition form factors and $(h, r, w_\pm)$ the $B \to M_1M_2$ transition form factors.

The matrix elements of $0 \to BB'$ are expressed as [27]

$$\langle BB'|q\gamma_\mu q'|0\rangle = \bar{u} F_{1,2} \gamma_\mu + \frac{F_2}{m_B} i \sigma_{\mu\nu} q^\nu.$$
with \( f_{M^0} \) the decay constant. For the \( B \to B^\prime \) transitions we have [13,26]

\[
\langle B^\prime | \bar{q} \gamma_\mu b | B \rangle = i \bar{u} [ i \gamma_\mu + g_2 \sigma_{\mu \nu} \tilde{p}^\nu + g_3 \tilde{p}_\mu + g_4 (p_B - p_B)^\mu ] \bar{y} v ,
\]

\[
\langle B^\prime | \bar{q} \gamma_\mu y b | B \rangle = i \bar{u} [ i f_1 \gamma_\mu + f_2 \sigma_{\mu \nu} \tilde{p}^\nu + f_3 \tilde{p}_\mu \\
+ f_4 (p_B - p_B)^\mu + f_5 (p_B - p_B)_\mu ] v ,
\]

\[
\langle B^\prime | \bar{q} b | B \rangle = i \bar{u} [ i \tilde{f}_1 \tilde{p} + \tilde{f}_2 (E_B - E_B) + \tilde{f}_3 (E_B - E_B)^\mu ] v ,
\]

\[
\langle B^\prime | \bar{q} y b | B \rangle = i \bar{u} [ i \tilde{f}_1 \tilde{p} + \tilde{f}_2 (E_B - E_B) + \tilde{f}_3 (E_B - E_B)^\mu ] v ,
\]

(11)

where \( \tilde{p}_\mu = (p_B - p_B)^\mu \), \( \tilde{g}_i (f_i) \) \( (j = 1, 2, 3) \) are the \( B \to B^\prime \) transition form factors.

The mesonic and baryonic form factors have momentum dependencies. For \( B \to M^{(*)} \), they are given by [32]

\[
F_A(q^2) = \frac{F_A(0)}{1 - \frac{q^2}{M_A^2}} \left[ 1 - \frac{\sigma_1 q^2}{M_A^2} + \frac{\sigma_2 q^2}{M_A^4} \right],
\]

\[
F_B(q^2) = \frac{F_B(0)}{1 - \frac{q^2}{M_B^2} + \frac{\sigma_1 q^2}{M_B^4}} ,
\]

(12)

where \( F_A = (F_{1BM}, V_1, A_0) \) and \( F_B = (F_{0BM}, A_{1,2}) \).

According to the approach of perturbative QCD counting rules, one presents the momentum dependencies of the form factors for \( B \to B^\prime, 0 \to B^\prime \) and \( B \to M_1 M_2 \) as [13,26,33–37]

\[
F_1 = \frac{\tilde{C}_f}{t^2}, \quad g_A = \frac{\tilde{C}_g}{t^2}, \quad f_S = \frac{\tilde{C}_S}{t^2}, \quad g_P = \frac{\tilde{C}_P}{t^2} ,
\]

\[
f_i = \frac{D_i}{t^2}, \quad g_i = \frac{D_i}{t^2}, \quad \tilde{f}_i = \frac{\tilde{D}_i}{t^2}, \quad \tilde{g}_i = \frac{\tilde{D}_i}{t^2} ,
\]

\[
h = \frac{C_h}{s^2}, \quad w_- = \frac{D_w}{s^2} ,
\]

(13)

where \( t \equiv (p_B + p_B^\prime)^2, s \equiv (p_1 + p_2)^2, \) and \( \tilde{C}_f = C_f (1/n(t/A_0^2))^{-\gamma} \) with \( \gamma = 2.148 \) and \( A_0 = 0.3 \) GeV. In Ref. [38], \( F_2 = F_1/(t \ln(t/A_0^2)) \) is calculated to be much less than \( F_1 \); hence we neglect it. Since \( h_A \) corresponds to the smallness of \( B(\bar{B}^0 \to p \bar{p}) \sim 10^{-8} \) [39–41], we neglect \( h_A \) as well. The terms \( (r, w_+) \) in Eq. (8) are neglected – following Refs. [36,37] – due to the fact that their parity quantum numbers disagree with the experimental evidence of \( J^P = 1^- \) for the meson-pair production [42].

The constants \( C_i (D_i) \) can be decomposed into sets of parameters that obey the SU(3) flavor and SU(2) spin symmetries. In Refs. [13,15,17,20,26,27,33,34,43,44], they are derived as

Table 1 The \( B^0 \to M^{(*)0} \) transition form factors at zero-momentum transfer, with \( (M_A, M_B) = (5.32, 5.32) \) and \( (5.27, 5.32) \) GeV for \( \pi^0 \) and \( \rho^0 \), respectively

| \( B^0 \to \pi^0, \rho^0 \) | \( F_{1 BM}^{\pi^0} \) | \( F_{0 BM}^{\pi^0} \) | \( V_1 \) | \( A_0 \) | \( A_1 \) | \( A_2 \) |
|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|
| \( \sqrt{2} f(0) \) | 0.29 | 0.29 | 0.31 | 0.30 | 0.26 | 0.24 |
| \( \sigma_1 \) | 0.48 | 0.76 | 0.59 | 0.54 | 0.73 | 1.40 |
| \( \sigma_2 \) | – | 0.28 | – | 0.10 | 0.50 | – |

\( C_{F_1}, C_{g_A} \)

\( = \frac{1}{9} (5C_{||} + 5C_{||} - C_{||}^2), \) \( (\text{for } \langle p \bar{p}(u\bar{u})_{V,A}(0)\rangle) \)

\( C_{F_1, C_{g_A}, C_{f_S}, C_{g_P}} \)

\( = \frac{1}{9} (C_{||} + 2C_{||}^* - 2C_{||}^* \tilde{C}_||, \tilde{C}_||^*), \) \( (\text{for } \langle p \bar{p}(d\bar{d})_{V,A,S,\rho}(0)\rangle) \)

\( D_{g_1, f_1}, D_{g_2, f_2} \)

\( = \frac{1}{5} (D_{||} - D_{||}^*, \tilde{D}_{||} - \tilde{D}_{||}^*) , \) \( (\text{for } \langle p \bar{p}(d\bar{d})_{V,A}(0)\rangle) \)

(14)

with \( j = 2, \ldots, 4, 5, C_{||}^* \equiv C_{||}^* + \delta C_{||} \text{ and } \tilde{C}_|| \equiv \tilde{C}_|| + \delta \tilde{C}_|| \). The direct CP violating asymmetry is defined as

\[
A_{CP}(B \to B^\prime X_M) \equiv \frac{\Gamma(B \to B^\prime X_M) - \Gamma(\tilde{B} \to \tilde{B}^\prime \bar{X}_M)}{\Gamma(B \to B^\prime X_M) + \Gamma(\tilde{B} \to \tilde{B}^\prime \bar{X}_M)} ,
\]

(15)

where \( B \to B^\prime \bar{X}_M \) denotes the anti-particle decay.

3 Numerical analysis

We use the following values for the numerical analysis. The CKM matrix elements are calculated via the Wolfenstein parameterization [1], with the world-average values

\[
\lambda = 0.22453 \pm 0.00044, \quad A = 0.836 \pm 0.015 , \\
\tilde{\rho} = 0.122^{+0.018}_{-0.017}, \quad \tilde{\eta} = 0.355^{+0.012}_{-0.011} .
\]

(16)

The decay constants are \( f_{1\nu}, f_{1\mu} \) \( = (30.4 \pm 0.2, 210.6 \pm 0.4) \) MeV [1], with \( (f_{1\nu}, f_{1\mu}) = (f_{1\nu}, f_{1\mu})/\sqrt{2} \). We adopt the \( B \to M^{(*)} \) transition form factors in Ref. [32], listed in Table 1. In Sect. 2, \( N_c \) has been presented as the phenomenological parameter determined by data. Empirically, one is able to determine \( N_c \) between 2 and \( \infty \). With the nearly universal value for \( N_c \) in the specific decays, the
factorization is demonstrated to be valid. For the tree-level internal W-emission dominated b-hadron decays, the extraction has given \( N_c \approx 2 \) that corresponds to \( a_2 \sim O(0.2 - 0.3) \) [15, 20, 45–49], where \( N_c \) differs due to the experimental uncertainties. For example, one obtains \( N_c = 2.15 \pm 0.17 \) in \( \Lambda_b \to B M_c \) [47, 48]. Here, we test if \( N_c \approx 2 \) can be used to explain the measured \( B(B^0 \to p \bar{p} \pi^0, p \bar{p} \pi^+ \pi^-) \).

The \( C_{h,w} \) for \( B^0 \to \pi^+ \pi^- \) and \( C_{l}(D_l) \) for \( 0 \to p \bar{p} \) (\( B^0 \to p \bar{p} \)) have been determined to be [15, 17, 20]

\[
(C_h, C_{w\ldots}) = (3.6 \pm 0.3, 0.7 \pm 0.2) \text{ GeV}^3,
(C_{l}, C_{\bar{l}}, C_{\bar{\pi}l}) = (154.4 \pm 12.1, 18.1
\pm 72.2, 537.6 \pm 28.7) \text{ GeV}^4,
(\delta C_{l}, \delta C_{l}', \delta C_{\bar{\pi}l}) = (19.3 \pm 21.6, -477.4 \pm 99.0, -342.3 \pm 61.4) \text{ GeV}^4,
(D_{l}, D_{\bar{l}}) = (45.7 \pm 33.8, -298.2 \pm 34.0) \text{ GeV}^5,
(D_{l}^2, D_{\bar{l}}^2, D_{\bar{l}}' D_{\bar{l}}') = (33.1 \pm 30.7, -203.6 \pm 133.4, 6.5
\pm 18.1, -147.1 \pm 29.3) \text{ GeV}^4,
(D_{\bar{l}}, D_{\bar{l}}', D_{\bar{l}}^2, D_{\bar{l}}') = (35.2 \pm 4.8, -38.2 \pm 7.5, -22.3 \pm 10.2, 504.5
\pm 32.4) \text{ GeV}^4.
\]

(17)

For \( a_2 \) in Eq. (7), the effective Wilson coefficients \( c_i^{eff} \) are calculated at the \( m_b \) scale in the NDR scheme, see Ref. [25]. They are related to the size of the decay, where the strong phases, together with the weak phase in \( V_{ub} \) and \( V_{td} \), play the key role in \( A_{CP} \).

Our results for the branching fractions and CP violating asymmetries of \( B^0 \to p \bar{p} X_M \) decays are summarized in Table 2, where we have averaged the particle and antiparticle contributions for the total branching fractions.

| Table 2 Decay branching fractions and direct CP asymmetries of \( B^0 \to p \bar{p} X_M \), where the first errors come from the estimations of the non-factorizable effects, the second ones from the uncertainties of the CKM matrix elements, and the third ones from those of the decay constants and form factors |

| \( 10^7 B(B^0 \to p \bar{p} \pi^0) \) | \( 5.0 \pm 1.9 \pm 0.3 \pm 0.9 \) | \( 5.0 \pm 1.9 [22] \) |
| \( 10^7 B(B^0 \to p \bar{p} \rho^0) \) | \( 1.8 \pm 1.7 \pm 0.1 \pm 0.4 \) | \( - \) |
| \( 10^8 B(B^0 \to p \bar{p} \pi^+ \pi^-) \) | \( 2.7 \pm 0.2 \pm 0.2 \pm 0.7 \) | \( 2.7 \pm 0.2 [19] \) |
| \( A_{CP}(B^0 \to p \bar{p} \pi^0) \) | \( (-16.8 \pm 4.8 \pm 1.6 \pm 1.8)\% \) | \( - \) |
| \( A_{CP}(B^0 \to p \bar{p} \rho^0) \) | \( (-12.6 \pm 2.2 \pm 1.2 \pm 1.7)\% \) | \( - \) |
| \( A_{CP}(B^0 \to p \bar{p} \pi^+ \pi^-) \) | \( (-11.4 \pm 0.2 \pm 1.2 \pm 1.4)\% \) | \( - \) |

4 Discussions and conclusions

The improved theoretical approaches such as QCD factorization (QCDF) and soft-collinear effective theory have been applied to two-body mesonic \( B \) decays [50–52]. Hence, the non-factorizable corrections of order \( 1/N_c^n \) with \( n = 1, 2 \) have been considered by calculating the vertex corrections from the hard gluon exchange and the hard spectator scattering. Unfortunately, there exist no similar approaches well applied to the \( B \to M_1 M_2 M_3, BB'M \) and \( BB'M' \) decays, due to the wave functions of \( B \to BB'(M'M') \) not as clear as those of \( B \to M \). By varying \( N_c \) from 2 to \( \infty \), one can still estimate the non-factorizable QCD effects with the corrections of order \( 1/N_c \). This relies on the generalized factorization, demonstrated to work well in \( B \to M_1 M_2 M_3, B \to BB', B \to BB'M (BB'M') \), \( B \to D \pi \) and \( \Delta b \to BM (A_+ \pi^-) \). We determine \( N_c = (2.15 \pm 0.20, 1.90 \pm 0.03) \) to interpret \( B(B^0 \to p \bar{p} \pi^0, p \bar{p} \pi^+ \pi^-) \) with \( \Delta N_c \) receiving the experimental uncertainties, which are indeed close to \( N_c \approx 2 \) used in \( B \to BB'M \) and \( \Delta b \to BM_{(c)} \) [15, 47–49].

In Table 2, \( B(B^0 \to p \bar{p} \pi^0) = 5.0 \times 10^{-7} \) receives the contributions from \( A_1, A_2 \) and their interference, denoted by \( A_{1 \times 2} \), which give \( B(B^0 \to p \bar{p} \pi^0) = B_1 + B_2 + B_{1 \times 2} \) with \( (B_1, B_2, B_{1 \times 2}) = (3.82, 0.33, 0.85) \times 10^{-7} \). The \( B_{1 \times 2} > 0 \) indicates constructive interference between \( A_1, A_2 \). By adopting \( N_c \) from \( B^0 \to p \bar{p} \pi^0 \), we predict \( B(B^0 \to p \bar{p} \rho^0) \). We find \( B(B^0 \to p \bar{p} \rho^0) \approx B(B^0 \to p \bar{p} \pi^0)/3 \) with \( (B_1, B_2, B_{1 \times 2}) = (2.00, 0.04, -0.24) \times 10^{-7} \). The minus sign of \( B_{1 \times 2} \) indicates destructive interference.

With the theoretical approach reasonably well established for the branching fractions, one can have reliable predictions for \( CP \) violation. For example, \( A_{CP}(B^0 \to p \bar{p} M^{(*)}) \) with \( M^{(*)} = (K^{*}, K^-, \pi^-) \) were predicted as \( (22 \pm 4, 6 \pm 1, -6 \pm 1)\% \) [43, 44], agreeing with the experimental values of \( (21 \pm 16, 2.1 \pm 2.0 \pm 0.4, -4.1 \pm 3.9 \pm 0.5)\% \) [1, 5]. Here, our predictions for \( A_{CP}(B^0 \to p \bar{p} \pi^0(\rho^0), p \bar{p} \pi^+ \pi^-) \) are around \(- (10 - 20)\% \). With \( \delta A_{CP} \) denoting the uncertainty for \( A_{CP} \), we present \( \delta A_{CP} \approx (0.2 - 0.3) A_{CP} \), which receives the theoretical uncertainties from the non-factorizable strong interaction, CKM matrix elements, form factors and decay constants.

Expressing the decay amplitude as \( A = T e^{i \delta_w} + P e^{i \delta_s} \), the CP asymmetry can be derived as

\[
A_{CP} = \frac{2 R \sin \delta_w \sin \delta_S}{1 + 2 R \cos \delta_w \cos \delta_S + R^2}.
\]

(18)

where \( \delta_w \) and \( \delta_S \) are the weak and strong phases arising from the tree (\( T \)) and penguin (\( P \))-level contributions, and
the ratio $R \equiv P/T$ suggests that a more suppressed $T$ amplitude is able to cause a more sizeable $\mathcal{A}_C$. Although $\bar{B}^0 \to p\bar{p}X_M$ involves complicated amplitudes, the relation in Eq. (18) can be used as a simple description for $\mathcal{A}_C(\bar{B}^0 \to p\bar{p}X_M)$. Being external and internal $W$-emission decays, $\bar{B}^- \to p\bar{p}\pi^-$ and $\bar{B}^0 \to p\bar{p}\pi^0$ proceed with $a_1 \sim O(1.0)$ and $a_2 \sim O(0.2 - 0.3)$ in the tree-level amplitudes [43,44], respectively. Consequently, the more suppressed $T$ amplitude with $a_2$ causes more interfering effect with the penguin diagrams, which corresponds to $|\mathcal{A}_C(B^0 \to p\bar{p}\pi^0)| > |\mathcal{A}_C(B^- \to p\bar{p}\pi^-)|$. In fact, we predict $|\mathcal{A}_C(B^0 \to p\bar{p}\pi^0)| = (16.8 \pm 5.4\%)$, which is three times larger than $|\mathcal{A}_C(B^- \to p\bar{p}\pi^-)|$ [43,44]. For the same reason, $|\mathcal{A}_C(B^0 \to p\bar{p}\rho^0, p\bar{p}\pi^+\pi^-)|$ can be as large as $(10 - 20\%)$.

Since $\mathcal{B}(\bar{B}^0 \to p\bar{p}\pi^0, p\bar{p}\pi^+\pi^-)$ are measured as large as $10^{-6}$, and well explained by the theory, with the predicted $|\mathcal{A}_C| > 10\%$, they become promising decays for measuring CP violation. By contrast, $\bar{B}^0 \to p\bar{p}\rho^0$ as well as the internal $W$-emission dominated $\Delta_b$ decays of $\Lambda_b^0 \to n\pi^0, n\rho^0$ have $\mathcal{B} \lesssim (1 - 2) \times 10^{-7}$, which make CP measurements a challenge even in the case of large $|\mathcal{A}_C| > 10\%$ [49].

In summary, we have investigated the branching fractions and direct CP violating asymmetries of the $\bar{B}^0 \to p\bar{p}\pi^0(\rho^0)$ and $\bar{B}^0 \to p\bar{p}\pi^+\pi^-$ decays. We have shown that these baryonic $B$-meson decays mediated dominantly through internal $W$-emission processes are promising processes to observe for the first time the CP violating effects in $B$ decays to final states with half-spin particles.

With a large predicted CP asymmetry $\mathcal{A}_C = (-16.8 \pm 5.4\%)$, which is accessible to the Belle II experiment, $\bar{B}^0 \to p\bar{p}\pi^0$ is particularly suited for a potential first observation of CP violation in baryonic $B$ decays in the coming years. Furthermore, the $\bar{B}^0 \to p\bar{p}\pi^+\pi^-$ decay, with its branching fraction of order $10^{-6}$ and the large predicted direct CP asymmetry $\mathcal{A}_C \sim -(10 - 20\%)$, is also in the realm of both Belle II and LHCb experiments.

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