Space-like Penguin Effects in $B \rightarrow PP$ Decays

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Abstract

The space-like penguin contributions to branching ratios and CP asymmetries in charmless decays of $B$ to two pseudoscalar mesons are studied using the next-to-leading order low energy effective Hamiltonian. Both the gluonic penguin and the electroweak penguin diagrams are considered. We find that the effects are significant.
1. Introduction

Penguin diagrams play an important role in charmless B-decays and direct CP violation [1,2]. But only time-like penguin diagrams were considered in the literature because they can provide the necessary different strong phases for CP violation by different loop effects of the internal $u$ and $c$ quarks [1]. The contribution of space-like penguin diagrams is usually neglected assuming form factor suppression. This assumption for neglecting space-like penguin effects is used not only for gluonic penguins but also for electro-weak penguins [3]. But it does not lie on a solid ground because the space-like penguin amplitudes can be remarkably enhanced by the hadronic matrix elements involving $(V-A)(V+A)$ or $(S+P)(S-P)$ currents [4]. Although space-like penguin diagrams can only provide an overall CP conserving phase due to final state interaction, it affects CP asymmetry by modifying the dispersive or absorptive parts of time-like penguin amplitudes. Obviously, it affects branching fractions too. In our recent paper [5], we illustrated the space-like penguin effects in CP asymmetries for the exclusive B-decays $B_u^- \rightarrow \bar{K}^0\pi^-$ and $K^0K^-$ using leading order Hamiltonian. In contrast to the naive expectation, the space-like penguin effects on CP asymmetries are found to be significant. In this letter we study space-like penguin effects in B to two pseudoscalar decays systematically. We concentrate on the charmless B decays because penguins play an important role in these decays. We use the next to leading order low energy effective Hamiltonian in order to consider both gluonic and electro-weak penguins. We can see later that the contribution of the electro-weak penguin is not negligible. The article is organized as following: In section II, we present the next to leading order effective Hamiltonian and the computation method. Section III devotes to the numerical results and corresponding discussions.
2. Effective hamiltonian and factorization approximation

Following ref. [6], the next-to-leading order low energy effective Hamiltonian describing \( \Delta B = -1, \Delta C = \Delta U = 0 \) transitions is given at the renormalization scale \( \mu = O(m_b) \) as

\[
H_{\text{eff}}(\Delta B = -1) = \frac{G_F}{\sqrt{2}} \left[ \sum_{q=u, c} v_q \left\{ Q_1^q C_1(\mu) + Q_2^q C_2(\mu) + \sum_{k=3}^{10} Q_k C_k(\mu) \right\} \right],
\]  

(1)

where the Wilson coefficient functions \( C_i(\mu) \) \( (i=1, \cdots, 10) \) are calculated in renormalization group improved perturbation theory and include leading and next-to-leading order QCD corrections and leading order corrections in \( \alpha \). The CKM factors \( v_q \) are defined as

\[
v_q = \begin{cases} 
V_{qd}^* V_{qb} & \text{for } b \to d \text{ transitions} \\
V_{qs}^* V_{qb} & \text{for } b \to s \text{ transitions.}
\end{cases}
\]

(2)

Here, we make use of the Wolfenstein parametrization[7] in which the CKM matrix can be written in terms of four parameters \( \lambda, A, \rho \) and \( \eta \) in the following form:

\[
V = \begin{pmatrix}
1 - \frac{1}{2} \lambda^2 & \lambda & \lambda^3 A(\rho - i \eta) \\
-\lambda & 1 - \frac{1}{2} \lambda^2 & \lambda^2 A \\
\lambda^3 A(1 - \rho - i \eta) & -\lambda^2 A & 1
\end{pmatrix}.
\]

(3)

The preferred values of the CKM Parameters are \( \lambda = 0.22, A = 0.8, \eta = 0.34, \rho = -0.12 \), which are obtained from the fit to experimental data[8]. The operators \( Q_1^u, Q_2^u, \)
$Q_3, \ldots, Q_{10}$ are given as the following forms:

$$Q_1^u = (\bar{q}_u u_\beta)_{V-A} (\bar{u}_\beta b_\alpha)_{V-A} \quad Q_2^u = (\bar{q}_u)_{V-A} (\bar{u} b)_{V-A}$$

$$Q_3 = (\bar{q} b)_{V-A} \sum_{q'} (\bar{q'} q')_{V-A} \quad Q_4 = (\bar{q}_a b_\beta)_{V-A} \sum_{q'} (\bar{q'} q'_\alpha)_{V-A}$$

$$Q_5 = (\bar{q} b)_{V-A} \sum_{q'} (\bar{q'} q')_{V+A} \quad Q_6 = (\bar{q}_a b_\beta)_{V-A} \sum_{q'} (\bar{q'} q'_\alpha)_{V+A}$$

$$Q_7 = \frac{3}{2} (\bar{q} b)_{V-A} \sum_{q'} e_{q'} (\bar{q'} q')_{V+A} \quad Q_8 = \frac{3}{2} (\bar{q}_a b_\beta)_{V-A} \sum_{q'} e_{q'} (\bar{q'} q'_\alpha)_{V+A}$$

$$Q_9 = \frac{3}{2} (\bar{q} b)_{V-A} \sum_{q'} e_{q'} (\bar{q'} q')_{V-A} \quad Q_{10} = \frac{3}{2} (\bar{q}_a q_\beta)_{V-A} \sum_{q'} e_{q'} (\bar{q'} q'_\alpha)_{V-A}.$$

where $Q_1^u$ and $Q_2^u$ are the current-current operators, and the current-current operators $Q_1^c$ and $Q_2^c$ are obtained from $Q_1^u$ and $Q_2^u$ through the substitution of $u \rightarrow c$. $Q_3, \ldots, Q_6$ are the QCD penguin operators, whereas $Q_7, \ldots, Q_{10}$ are the electroweak penguin operators. The quark $q = d$ or $s$ is for $b \rightarrow d$ or $s$ transitions, respectively; $q'$ is running over the quark flavours being active at the scale $\mu = O(m_b)$ ($q' \in \{u, d, c, s, b\}$); $e_{q'}$ are the corresponding quark charges; the indices $\alpha, \beta$ are $SU(3)_c$ color indices; $(V \pm A)$ refer to $\gamma \mu (1 \pm \gamma_5)$. It should be noted that the Hamiltonian (1) can be viewed as the generalization of the leading logarithmic Hamiltonians presented in [9,10].

Beyond the leading logarithmic approximation, the Wilson coefficient functions $C_i(\mu)$ depend both on the form of the operator basis (4) and on the renormalization scheme. Here, we use the renormalization scheme independent Wilson coefficient functions [11]:

$$C(\mu) = \left[ 1 + \frac{\alpha_s(\mu)}{4\pi} \hat{r}_s^T + \frac{\alpha(\mu)}{4\pi} \hat{r}_e^T \right] \cdot C(\mu),$$

where $\hat{r}_s$ and $\hat{r}_e$ are obtained from one-loop matching conditions. Now, taking the QCD and electroweak one-loop level matrix elements of the operators $Q_i$ ($Q_i = Q_1^u, Q_2^u, Q_3, \ldots, Q_{10}$) into account through

$$< Q^T(\mu) > = < Q^T >_0 \left[ 1 + \frac{\alpha_s(\mu)}{4\pi} \hat{m}_s^T(\mu) + \frac{\alpha_{em}}{4\pi} \hat{m}_e^T(\mu) \right],$$

4
which define matrices $\hat{m}_s(\mu)$ and $\hat{m}_e(\mu)$. In Eq. (5) and (6), $C(\mu)$, $\bar{C}(\mu)$ and $Q$ are all column vectors, where the vector $Q$ are given by the operator basis $Q_i$, and $<Q>^0$ denote the tree level matrix elements of these operators. Combine Eq.(5) and (6), we obtain

$$<Q^T(\mu) \cdot C(\mu)>\equiv <Q^T>_0 \cdot C'(\mu)$$

where $C'(\mu)$ are defined as

$$C'_1 = \bar{C}_1, \quad C'_2 = \bar{C}_2,$$

$$C'_3 = \bar{C}_3 - P_s/3, \quad C'_4 = \bar{C}_4 + P_s,$$

$$C'_5 = \bar{C}_5 - P_s/3, \quad C'_6 = \bar{C}_6 + P_s,$$

$$C'_7 = \bar{C}_7 + P_e, \quad C'_8 = \bar{C}_8,$$

$$C'_9 = \bar{C}_9 + P_e, \quad C'_{10} = \bar{C}_{10},$$

where $P_{s,e}$ are given by

$$P_s = \frac{\alpha_s}{8\pi} \bar{C}_2(\mu) \left[ \frac{10}{9} - G(m, q, \mu) \right],$$

$$P_e = \frac{\alpha_{em}}{9\pi} \left( 3\bar{C}_1 + \bar{C}_2(\mu) \right) \left[ \frac{10}{9} - G(m, q, \mu) \right],$$

$$G(m, q, \mu) = -4 \int_0^1 dx \ x(1-x)ln \left[ \frac{m^2 - x(1-x)q^2}{\mu^2} \right],$$

here $q = u, c$, for numerical calculation, we take $m_u = 0.005 GeV, m_c = 1.35 GeV$, and $q^2$ denotes the momentum transfer squared of the virtual gluons, photons and $Z^0$ appearing in the QCD and electroweak penguin matrix elements. For the details of this calculation, the reader is referred to ref. [12,13].

The renormalization scheme independent Wilson coefficient functions $\bar{C}_i(\mu)$ at the scale $\mu = O(m_b)$ are obtained by first calculating the Wilson coefficients at $\mu = O(m_W)$ and then using the renormalization group equation to evolve them down to $O(m_b)$. We
use in our analysis, $\alpha_s(m_Z) = 0.118$, $\alpha(m_Z) = 1/128[14]$ and $m_t = 174\text{GeV}[15]$ and the numerical values of the renormalization scheme independent Wilson Coefficients $\bar{C}_i(\mu)$ at $\mu = O(m_b)$ are[13]

$$
\begin{align*}
\bar{c}_1 &= -0.313, \quad \bar{c}_2 = 1.150, \quad \bar{c}_3 = 0.017, \\
\bar{c}_4 &= -0.037, \quad \bar{c}_5 = 0.010, \quad \bar{c}_6 = -0.046, \\
\bar{c}_7 &= -0.001 \cdot \alpha_{em}, \quad \bar{c}_8 = 0.049 \cdot \alpha_{em}, \quad \bar{c}_9 = -1.321 \cdot \alpha_{em}, \\
\bar{c}_{10} &= 0.267 \cdot \alpha_{em}.
\end{align*}
$$

(10)

With the help of Eq. (7), the two-body decay amplitude $< PP'|H_{eff}(\Delta B = -1)|B >$ can be expressed as linear combinations of $< PP'|Q_i|B >_0$. The hadronic matrix elements $< PP'|Q_i|B >_0$ are evaluated using the factorization approximation [16]. It should be noted that this approach has already been used in the literature to analyze the QCD or electroweak time-like penguin contributions[12]. However, we go further in this letter by including the space-like penguin diagrams. As in [5,12], we also neglect W-annihilation or W-exchange diagram contributions in our present analysis which are commonly assumed to be form factor suppressed.

Using the vacuum-saturation approximation, the decay amplitude $< PP'|H_{eff}|B >$ can be factorized into a product of two current matrix elements $< P|J^\mu|0 >$ and $< P'|J_\mu^\prime|B >$ for the tree and time-like penguin diagrams (Fig.1), or the product of $< pp'|J^\mu|0 >$ and $< 0|J_\mu^\prime|B >$ for the space-like penguin diagrams (Fig.2). In this work, the hadronic matrix elements are calculated in BSW method[16]. We define

$$
M_{q_1q_2q_3}^{pp'} = < P| (\bar{q}_1 q_2)_{V-A}|0 > < P'| (\bar{b}_3 q_3)_{V-A}|B >,
$$

(11)

and

$$
S_{q_1q_2q_3}^{pp'} = < PP'| (\bar{q}_1 q_2)_{V-A}|0 > < 0| (\bar{b}_3 q_3)_{V-A}|B >.
$$

(12)
where $M^{pp'}_{q1q2q3}$ denotes the hadronic matrix element in tree and time-like penguin diagram case, while $S^{pp'}_{q1q2q3}$ denotes space-like penguin case. When the (V-A)(V+A) current are transformed into (S+P)(S-P) and further into (V-A)(V-A) ones using equation of motion for the time-like and space-like penguin amplitudes, there appear the terms which are proportional to $\frac{2m_B^2}{(m_q-m_{q'})^2}$ and $\frac{2m_B^2}{(m_q-m_{q'})^2}$, respectively. If $q = q'$ as in the decay modes:

$$\bar{B}_d \to \pi^- \pi^+, \pi^0 \pi^0, \pi^0 \eta, \pi^0 \eta', \eta \eta', \eta \eta, K^0 \bar{K}^0; \quad \text{(for } b \to d \text{ transitions)}$$

$$\bar{B}_s \to K^- K^+, \bar{K}^0 K^0, \pi^0 \eta, \pi^0 \eta', \eta \eta', \eta \eta'. \quad \text{(for } b \to s \text{ transitions)},$$

the denominator of the factor $\frac{2m_B^2}{(m_q-m_{q'})^2}$ is zero. So, we can not use equation of motion to compute the amplitudes of these decays. We have to compute the matrix elements of (S+P)(S-P) operators directly. We shall discuss it elsewhere.

As an example of how to factorize the decay amplitudes into the product of hadronic matrix elements, we give the result of $<\pi^- \pi^0|H_{eff}|\bar{B}_u^->$ in the following,

$$\begin{align*}
<\pi^- \pi^0|H_{eff}|\bar{B}_u^-> &= \\
&= \frac{G_F}{\sqrt{2}} \sum_{q=u,c} v_q \left\{ a_1 \delta_{uq} + a_3 - \frac{2M^2_{\pi^-}}{(m_d+m_u)(m_u-m_b)} (a_5 + a_7) + a_9 \right\} M_{d\bar{u}u}^{\pi^- \pi^0} + \\
&+ \left\{ a_2 \delta_{uq} - a_3 + \frac{2M^2_{\pi^0}}{m_d(m_d-m_b)} (a_5 - a_7/2) - \frac{3}{2} a_8 + \frac{3}{2} a_{10} + \frac{1}{2} a_9 \right\} M_{d\bar{u}u}^{\pi^0 \pi^-} + \\
&+ \left\{ a_3 + \frac{2M^2_B}{(m_u+m_b)(m_d-m_u)} (a_5 + a_7) + a_9 \right\} (S_{d\bar{u}u}^{\pi^- \pi^0} + S_{d\bar{u}u}^{\pi^0 \pi^-}) \right], \quad \text{(13)}
\end{align*}$$

where the term $(S_{d\bar{u}u}^{\pi^- \pi^0} + S_{d\bar{u}u}^{\pi^0 \pi^-})$ is the contribution obtained from two space-like penguin diagrams, and the quark masses are taken as $m_d = 0.01GeV$, $m_u = 0.005GeV$, $m_s = 0.175GeV$, $m_b = 5.0GeV$. $a_k$ is defined as

$$a_{2i-1} \equiv \frac{C'_0}{3} + C'_2,$$

$$a_{2i} \equiv C'_{2i-1} + \frac{C'_0}{3}, \quad (i = 1, 2, 3, 4, 5)$$
The general expression for the one-body pseudoscalar matrix element of the axial-vector is

\[ \langle 0| V_\mu - A_\mu | P(q) \rangle = i f_P q_\mu, \]

(14)

where \( q \) represents the momentum of the pseudoscalar meson, and \( f_P \) is the decay constant. The two-body pseudoscalar-pseudoscalar matrix element of the vector current is

\[ \langle P_2(q_2)| V_\mu - A_\mu | P_1(q_1) \rangle = f_+(q_+^2)q_+ + f_-(q_-^2)q_- \]

(15)

where \( q_\pm = q_1 \pm q_2 \), and the form factor \( f_\pm \) is given by the monopole parametrization

\[ f_+(q_+^2) \approx \frac{f_+(0)}{1 - q^2/m_{\text{pole}}^2}, \]

(16a)

\[ f_-(q_-^2) = -\frac{m_1 - m_2}{m_1 + m_2} f_+(q_-^2). \]

(16b)

With equation (14~16), we obtain

\[ M_{\text{dau}}^{\pi^0 - \pi^0} = -\frac{i}{\sqrt{2}} f_\pi^\ast f_{B^0}^{\pi^0} (M_{B^+}^2 - M_{\pi^0}) \frac{M_{B^+} - M_{\pi^0}}{M_{B^+} + M_{\pi^0}} \left[ (M_{B^+} + M_{\pi^0})^2 - M_{\pi^-}^2 \right], \]

\[ M_{\text{uud}}^{\pi^0 \pi^-} = -if_{\pi^0}^\ast f_{B^-}^{\pi^-} (M_{B^+}^2 - M_{\pi^-}) \frac{M_{B^-} - M_{\pi^-}}{M_{B^-} + M_{\pi^-}} \left[ (M_{B^-} + M_{\pi^-})^2 - M_{\pi^0}^2 \right], \]

(17)

and

\[ S_{\text{dau}}^{\pi^0 - \pi^0} = -\frac{i}{\sqrt{2}} f_{B^-} f_{B^-} (M_{B^+}^2 - M_{\pi^0}) \frac{M_{B^-} - M_{\pi^0}}{M_{\pi^-} + M_{\pi^-}} \left[ (M_{\pi^-} + M_{\pi^0})^2 - M_{B^-}^2 \right], \]

\[ S_{\text{dau}}^{\pi^0 \pi^-} = -\frac{i}{\sqrt{2}} f_{B^-} f_{B^-} (M_{B^+}^2 - M_{\pi^-}) \frac{M_{\pi^-} - M_{\pi^0}}{M_{\pi^-} + M_{\pi^-}} \left[ (M_{\pi^-} + M_{\pi^0})^2 - M_{B^-}^2 \right], \]

(18)

where the factors \( \frac{i}{\sqrt{2}}, -\frac{i}{\sqrt{2}} \) come from the constituent of \( \pi^0 = \frac{1}{\sqrt{2}} (\bar{u}u - d\bar{d}) \).

In order to give numerical results, we need to know the form factors. For the decay form factors like \( f_+^{B\pi}(M^2) \), etc., we can use BSW [16] method to calculate them. For the annihilation form factor \( f_+^a(Q^2) \), we do not have reliable method to compute it. But at \( Q^2 = M_B^2 \) one is far from the \( K\pi, \pi\pi, \eta K \) resonance region. So, for the charmless \( B \) decays, because the large energy release, we can use the form factor in its asymptotic form. For charmless \( B \) to two pseudoscalars decays, the asymptotic form factor predicted
by QCD [18] should be a resonable approximation. So we take \( f_+^a(Q^2) = i16\pi\alpha_s f_B^2/Q^2 \).

Now we are in a position to give the numerical results.

### 3. Results and discussions

The decay width of a B-meson at rest decaying into two pseudoscalars is

\[
\Gamma(B \to PP') = \frac{1}{8\pi} |< PP'|H_{eff}|B>|^2 \frac{|p|}{M_B^2},
\]

where

\[
|p| = \frac{((M_B^2 - (M_P + M_{P'}))^2)(M_B^2 - (M_P - M_{P'}))^2)}{2M_B}^{1/2}
\]

is the momentum of the pseudoscalar meson \( P \) or \( P' \). The corresponding branching ratios are given by

\[
Br(B \to PP') = \frac{\Gamma(B \to PP')}{\Gamma_{tot}}.
\]

In our numerical calculation, we take \( \Gamma_{tot}^{B_u} = 4.27 \times 10^{-13}\text{GeV}, \Gamma_{tot}^{B_d} = 4.39 \times 10^{-13}\text{GeV}, \)
and \( \Gamma_{tot}^{B_s} = 4.91 \times 10^{-13}\text{GeV} \).

In order to obtain the CP-violating parameter, the B-meson decay amplitude can be generally expressed as

\[
< PP'|H_{eff}|B> = \frac{G_F}{\sqrt{2}} \sum_{q=u,c} v_q (C'_1 < Q_1^q > + C'_2 < Q_2^q > \\
+ \sum_{k=3}^{10} C'_k < Q_k >)
\equiv \frac{G_F}{\sqrt{2}} \sum_{q=u,c} v_q F_q.
\]

With the help of eq.(22), one can get the CP-violating asymmetry parameter

\[
\mathcal{A}_{cp} = \frac{\Gamma(B \to PP') - \Gamma(B \to PP')}{\Gamma(B \to PP') + \Gamma(B \to PP')} = \frac{2Im(v_q v_q^*) Im(F_c^*)}{v_u^2 + v_c^2 (F_u^*)^2 + 2Re(F_c^*)}.
\]
Since the branching ratios and CP asymmetries depend crucially on the parameter $q^2$ describing the momentum squared of the exchanged virtual particles appearing in the penguin matrix elements of Fig. 1 and 2, we should consider it in detail. Here, we use the same simple picture for two-body decays illustrated in Fig. 1c and 2 as the one in ref [5]. With the simple physical picture presented in Ref.[5], the average value of the momentum squared $<q^2>$ of the exchanged virtual particles can be given by

$$<q^2> = m_b^2 + m_q^2 - 2m_bE_q,$$  \hspace{1cm} (24)

where $E_q$ is determined from

$$E_q + \sqrt{E_q^2 - m^2_q + m^2_q'} + \sqrt{4E_q^2 - 4m^2_q + m^2_q'} = m_b; \hspace{1cm} (25a)$$

for the time-like penguin channels; or from

$$E_q + \sqrt{E_q^2 - m^2_q + m^2_q'} = m_b + m_q'; \hspace{1cm} (25b)$$

for the space-like penguin channels. When we factorize $<Q_k>_0$ of hairpin diagrams illustrated in Fig. 1d), we find $<Q_3>_0 = -<Q_5>_0$, $<Q_4>_0 = -<Q_6>_0$, and $<Q_7>_0 = -<Q_9>_0$. and hence the factor in Eq. 6:

$$\left\{ \frac{\alpha_s}{8\pi} \left[ -\frac{1}{N_c} <Q_3>_0 + <Q_4>_0 - \frac{1}{N_c} <Q_5>_0 + <Q_6>_0 \right] \tilde{C}_2(\mu) \\
+ \frac{\alpha}{3\pi} \left[ <Q_7>_0 + <Q_9>_0 \right] \left[ \bar{C}_1(\mu) + \frac{1}{N_c} \bar{C}_2(\mu) \right] \right\}$$  \hspace{1cm} (26)

vanishes because of the cancelation. So, we do not need to consider the hairpin diagrams.

The numerical results of the space-like penguin contributions to the branching ratios and CP-violating asymmetries are given in table 1. In the meantime, we also calculate the branching ratios and CP-violating asymmetries with only the tree and time-like penguin contributions for comparison. We also present the results with only tree and gluonic penguin contributions. All the parameters such as meson decay constants and form factors.
needed in our calculation are taken as $f_{\pi^\pm} = 0.13 GeV$, $f_K = 0.160 GeV[14]$, $f_{\pi^0}^{\bar uu} = -f_{\pi^0}^{\bar dd} = f_{\pi^\pm}/\sqrt{2}$, $f_{\eta}^{\bar uu} = f_{\eta}^{\bar dd} = 0.092$, $f_{\eta}^{\bar ss} = -0.105$, $f_{\eta'}^{\bar uu} = f_{\eta'}^{\bar dd} = 0.049$, $f_{\eta'}^{\bar ss} = 0.12[4]$, $f_D = 0.23$, $f_{D_s} = 0.281[19]$, $f_B = 1.5 \times f_{\pi^\pm}[20]$, $f_{B_s} = 0.206[21]$, and $f_B^{B_u \pi^-}(0) = 0.29$, $f_B^{B_u K^-}(0) = 0.32[22]$, $f_B^{B_u \eta_{uu}}(0) = 0.307$, $f_B^{B_u \eta_{uu}'}(0) = 0.254$, $f_B^{BD}(0) = 0.690[16]$, $f_B^{B_u \eta_{ss}}(0) = 0.335$, $f_B^{B_u \eta_{ss}'}(0) = 0.282[23]$, $f_B^{B_s D}(0) = 0.648[24]$.

From Table 1 we can see the following features:

(i) For most of the charmless decays, penguin contributions are important.

(ii) For $B_u^- \to \pi^- \eta$, $K^- \pi^0$, $K^- \eta$, $K^- \eta'$, $\bar B_d^0 \to \bar K^0 \pi^0$, $\bar K^0 \eta$, $\bar K^0 \eta'$ and $B_s^0 \to \pi^0 K^0$, the contribution of the electro-weak penguins are not negligible.

(iii) The space-like penguin effects in $B_u^- \to \pi^- \pi^0$ are amazingly large. The correction to the branching ratio is more than 100%, while to the CP asymmetry is about an order of magnitude, actually, $A_{cp} \sim 0.77\%$ with only time-like penguin, but $A_{cp} \sim -70.4\%$ when including the space-like penguin.

For $B_u^- \to K^- \pi^0$, $K^- \eta$, $K^- \eta'$, $K^0 \pi^-$, $\bar B_d^0 \to K^- \pi^+$, $\bar K^0 \eta$, $\bar K^0 \eta'$, and $B_s^0 \to \pi^- K^+$, $\pi^0 K^0$, the space-like penguin contributions are also dominant.

In $B_u^- \to \pi^- \eta$, $\pi^- \eta'$, the space-like penguin contribution is zero. The reason is that there are two space-like penguin diagrams in each channel and the contributions of the two diagrams exactly cancel each other.

In general, we can conclude that the space-like penguin effects are not negligible in most of the charmless two-pseudoscalar decays of the B mesons. The space-like penguins can affect not only CP asymmetries, but also decay branching ratios. Further investigations are definitely needed.

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Table 1. The branching ratios and the CP asymmetries, where the “Only Tree” means the branching ratios with only tree diagram contribution, “T-like” denotes the time-like penguin contributions, the “S-like” denotes the space-like penguin contributions, “QCD” means QCD penguin and tree diagrams contributions, and “QCD+EW” denotes full tree, QCD and EW(electro-weak) penguin contributions.
| decay mode               | Only Tree | Tree+T-like | Tree+T-like+S-like | Tree+T-like |
|--------------------------|-----------|-------------|-------------------|-------------|
|                          | QCD       | QCD+EW     | QCD               | QCD+EW     |
| $B_u^- \to \pi^- \pi^0$  | $2.96 \times 10^{-6}$ | $2.90 \times 10^{-6}$ | $2.74 \times 10^{-6}$ | $1.30 \times 10^{-6}$ | $1.12 \times 10^{-6}$ | $0.69\%$ | $0.77\%$ |
| $B_u^- \to \pi^- \eta$  | $2.19 \times 10^{-6}$ | $3.68 \times 10^{-6}$ | $2.81 \times 10^{-6}$ | $3.68 \times 10^{-6}$ | $2.81 \times 10^{-6}$ | $35.1\%$ | $33.5\%$ |
| $B_u^- \to \pi^- \eta'$ | $6.95 \times 10^{-7}$ | $5.84 \times 10^{-6}$ | $5.8 \times 10^{-6}$ | $5.84 \times 10^{-6}$ | $5.8 \times 10^{-6}$ | $16.0\%$ | $16.1\%$ |
| $B_u^- \to K^0 K^-$     | $0$       | $4.16 \times 10^{-7}$ | $4.08 \times 10^{-7}$ | $3.79 \times 10^{-7}$ | $3.72 \times 10^{-7}$ | $2.47\%$ | $2.49\%$ |
| $B_u^- \to K^- \pi^0$   | $2.12 \times 10^{-7}$ | $2.93 \times 10^{-6}$ | $4.22 \times 10^{-6}$ | $5.29 \times 10^{-6}$ | $6.55 \times 10^{-6}$ | $-8.58\%$ | $-6.15\%$ |
| $B_u^- \to K^- \eta$    | $1.57 \times 10^{-7}$ | $1.53 \times 10^{-7}$ | $1.84 \times 10^{-7}$ | $4.00 \times 10^{-8}$ | $7.13 \times 10^{-8}$ | $6.17\%$ | $4.38\%$ |
| $B_u^- \to K^- \eta'$   | $4.98 \times 10^{-8}$ | $8.19 \times 10^{-6}$ | $7.74 \times 10^{-6}$ | $8.03 \times 10^{-6}$ | $7.59 \times 10^{-6}$ | $-3.35\%$ | $-3.53\%$ |
| $B_u^- \to K^0 \pi^-$   | $0$       | $4.95 \times 10^{-6}$ | $4.86 \times 10^{-6}$ | $7.10 \times 10^{-6}$ | $6.97 \times 10^{-6}$ | $-0.18\%$ | $-0.18\%$ |
| $\bar{B}_d^0 \to K^- \pi^+$ | $3.77 \times 10^{-7}$ | $5.78 \times 10^{-6}$ | $5.96 \times 10^{-6}$ | $1.04 \times 10^{-5}$ | $1.06 \times 10^{-5}$ | $-8.23\%$ | $-8.0\%$ |
| $\bar{B}_d^0 \to K^0 \pi^0$ | $6.76 \times 10^{-10}$ | $2.46 \times 10^{-6}$ | $1.56 \times 10^{-6}$ | $3.64 \times 10^{-6}$ | $2.74 \times 10^{-6}$ | $0.44\%$ | $0.73\%$ |
| $\bar{B}_d^0 \to K^0 \eta$ | $6.72 \times 10^{-10}$ | $4.25 \times 10^{-8}$ | $1.97 \times 10^{-9}$ | $6.42 \times 10^{-8}$ | $2.33 \times 10^{-8}$ | $1.40\%$ | $47.0\%$ |
| $\bar{B}_d^0 \to K^0 \eta'$ | $1.87 \times 10^{-10}$ | $7.97 \times 10^{-6}$ | $7.42 \times 10^{-6}$ | $7.64 \times 10^{-6}$ | $7.09 \times 10^{-6}$ | $-0.42\%$ | $-0.44\%$ |
| $\bar{B}_s^0 \to \pi^- K^+$ | $3.51 \times 10^{-6}$ | $3.08 \times 10^{-6}$ | $3.07 \times 10^{-6}$ | $1.70 \times 10^{-6}$ | $1.68 \times 10^{-6}$ | $8.34\%$ | $8.41\%$ |
| $\bar{B}_s^0 \to \pi^0 K^0$ | $1.13 \times 10^{-8}$ | $1.18 \times 10^{-7}$ | $7.06 \times 10^{-8}$ | $3.36 \times 10^{-7}$ | $2.94 \times 10^{-7}$ | $-6.05\%$ | $-10.2\%$ |
| $\bar{B}_s^0 \to \eta K^0$ | $1.12 \times 10^{-8}$ | $1.48 \times 10^{-6}$ | $1.62 \times 10^{-6}$ | $1.53 \times 10^{-6}$ | $1.67 \times 10^{-6}$ | $5.27\%$ | $4.86\%$ |
| $\bar{B}_s^0 \to \eta' K^0$ | $3.13 \times 10^{-9}$ | $8.33 \times 10^{-6}$ | $8.23 \times 10^{-6}$ | $8.14 \times 10^{-6}$ | $8.04 \times 10^{-6}$ | $3.09\%$ | $3.11\%$ |
**Figure Captions**

**Fig. 1.** Quark diagrams for a $B$ meson decaying into two light pseudoscalar mesons $P$ and $P'$ through the tree process $b \rightarrow u(\bar{u}q)$: a) the internal $W$-emission diagram, (b) the external $W$-emission diagram; and the time-like penguin diagram process $b \rightarrow q(q'q'')$: c) the time-like pure penguin diagram, and d) the time-like hairpin diagram. The subscripts “s” denote “spectator”. The dark dot stands for the contraction of the $W$-loop.

**Fig. 2.** Quark diagrams for a $B$ meson decaying into two light pseudoscalar mesons $P$ and $P'$ through the space-like penguin process $(bq') \rightarrow (qq')$. The subscripts “v” denote “vacuum”. The dark dot stands for the contraction of the $W$-loop.
