Radiative Baryonic $B$ Decays

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Abstract

We study the structure-dependent contributions to the radiative baryonic $B$ decays of $B \to B\bar{B}'\gamma$ in the standard model. We show that the decay branching ratios of $Br(B \to B\bar{B}'\gamma)$ are $O(10^{-7})$, which are larger than the estimated values of $O(10^{-9})$ induced from inner bremsstrahlung effects of the corresponding two-body modes. In particular, we find that $Br(B^- \to \Lambda\bar{p}\gamma)$ is around $1 \times 10^{-6}$, which is close to the pole model estimation but smaller than the experimental measurement from BELLE.

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The radiative baryonic $B$ decays of $B \to \mathbf{B}\mathbf{B}'\gamma$ are of interest since they are three-body decays with two spin-1/2 baryons ($\mathbf{B}$ and $\mathbf{B}'$) and one spin-1 photon in the final states. The rich spin structures allow us to explore various interesting observables such as triple momentum correlations to investigate CP or T violation [1, 2]. Moreover, since these radiative decays could dominantly arise from the short-distance electromagnetic penguin transition of $b \to s\gamma$ [3] which has been utilized to place significant constraints on physics beyond the Standard Model (SM) [4, 5], they then appear to be the potentially applicable probes to new physics.

There are two sources to produce radiative baryonic $B$ decays. One is the inner bremsstrahlung (IB) effect, in which the radiative baryonic $B$ decays of $B \to \mathbf{B}\mathbf{B}'\gamma$ are from their two-body decay counterparts of $B \to \mathbf{B}\mathbf{B}'$ via the supplementary emitting photon attaching to one of the final baryonic states. Clearly, the radiative decay rates due to the IB contributions are suppressed by $\alpha_{em}$ comparing with their counterparts. According to the existing upper bounds of $B \to \mathbf{B}\mathbf{B}'$, given by [6, 7, 8]

\begin{align}
Br(\bar{B}^0 \to p\bar{p}) &< 2.7 \times 10^{-7} \text{ (BABAR)}, \\
Br(\bar{B}^0 \to \Lambda\bar{\Lambda}) &< 7.9 \times 10^{-7} \text{ (BELLE)}, \\
Br(B^- \to \Lambda\bar{p}) &< 4.6 \times 10^{-7} \text{ (BELLE)},
\end{align}

one finds that

$$Br(B \to \mathbf{B}\mathbf{B}'\gamma)_{IB} \leq O(10^{-9}).$$

(2)

Unfortunately, the above branching ratios are far from the present accessibility at the B factories of BABAR and BELLE. However, the other source, which is the structure-dependent (SD), is expected to enhance the decays of $Br(B \to \mathbf{B}\mathbf{B}'\gamma)$, such as $B \to \Lambda\bar{p}\gamma$ arising from $b \to s\gamma$ [1, 9, 10]. With the large branching ratio of $b \to s\gamma$ [11, 12] in the range of $10^{-4}$ we expect that $Br(B^- \to \mathbf{B}\mathbf{B}'\gamma)$ could be as large as $Br(B^- \to \mathbf{B}\mathbf{B}')$. In this report, we shall concentrate on the SD contributions to $Br(B \to \mathbf{B}\mathbf{B}'\gamma)$.

To start our study, we must tackle the cumbersome transition matrix elements in $B \to \mathbf{B}\mathbf{B}'$. As more and more experimental data on three-body decays [13, 14, 15] in recent years, the theoretical progresses are improved to resolve the transition matrix element problems. One interesting approach is to use the pole model [16, 17] through the intermediated particles and another one is to rely on the QCD counting rules [18, 19, 20] by relating the transition
matrix elements with three form factors and fitting with experimental data. In Ref. [9], Cheng and Yang have worked out the radiative baryonic B decays based on the pole model. In this paper, we handle the transition matrix elements according to the QCD counting rules.

We begin with the decay of $B^- \rightarrow \Lambda \bar{p} \gamma$. As depicted in Fig. 1 in the SM the relevant Hamiltonian due to the SD contribution for $B^- \rightarrow \Lambda \bar{p} \gamma$ is

$$\mathcal{H}_{SD} = -\frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* c_7^{\text{eff}} O_7, \quad (3)$$

with the tensor operator

$$O_7 = \frac{e}{8\pi^2} m_b \bar{s} \sigma_{\mu\nu} F^{\mu\nu}(1 + \gamma_5) b, \quad (4)$$

where $V_{tb} V_{ts}^*$ and $c_7^{\text{eff}}$ are the CKM matrix elements and Wilson coefficient, respectively, and the decay amplitude is found to be

$$A(B^- \rightarrow \Lambda \bar{p} \gamma) = \frac{G_F}{\sqrt{2}} V_{ts} V_{tb}^* \frac{e}{8\pi^2} 2 c_7^{\text{eff}} \left\{ m_b^2 \epsilon^{\mu} \langle \Lambda \bar{p} | \bar{s} \gamma_{\mu} (1 - \gamma_5) b | B^- \rangle - 2 m_b p_B \cdot \epsilon \langle \Lambda \bar{p} | \bar{s} (1 + \gamma_5) b | B^- \rangle \right\}, \quad (5)$$
where we have used the condition \( m_b \gg m_s \) such that the terms relating to \( m_s \) are neglected. We note that Eq. (5) is still gauge invariant.

In order to solve the encountered transition matrix elements in Eq. (5), we write the most general form

\[
\langle \Lambda \bar{p}|\bar{s}\gamma_{\mu}b|B^-\rangle = i\bar{u}(p_{\Lambda})[a_1\gamma_{\mu}\gamma_5 + a_2p_{\mu}\gamma_5 + a_3(p_{\bar{p}} - p_{\Lambda})\gamma_5]v(p_{\bar{p}}),
\]

\[
\langle \Lambda \bar{p}|\bar{s}\gamma_{\mu}\gamma_5 b|B^-\rangle = i\bar{u}(p_{\Lambda})[c_1\gamma_{\mu} + c_2\sigma_{\mu\nu}p^\nu + \gamma_{\mu}(p_{\bar{p}} + p_{\Lambda})]v(p_{\bar{p}}),
\]

where \( p = p_B - p_{\Lambda} - p_{\bar{p}} \) and \( a_i(c_i) \) \((i = 1, \ldots, 3)\) are form factors.

To find out the coefficients \( a_i(c_i) \) in Eq. (5), we invoke the work of Chua, Hou and Tsai in Ref. [20]. In their analysis, three form factors \( F_A, F_P \) and \( F_V \) are used to describe \( B \to B\bar{B}' \) transitions based on the QCD counting rules [18], that require the form factors to behave as inverse powers of \( t = (p_B + p_{B'})^2 \). The detail discussions can be referred to Refs. [19, 20]. In this paper, we shall follow their approach. The representations of the matrix elements for the \( B^- \to p\bar{p} \) transition are given by [20]

\[
\langle p\bar{p}|\bar{s}(1 \pm \gamma_5)b|B^-\rangle = i\bar{u}(p_{\bar{p}})[(F_A \not{p}\gamma_5 \pm F_V \not{p}) + (F_P \gamma_5 \pm F_S)]v(p_{\bar{p}}),
\]

with a derived relation \( F_S = F_P \). In terms of the approach of [19, 20], those of the \( B^- \to \Lambda \bar{p} \) transition are given by

\[
\langle \Lambda \bar{p}|\bar{s}(1 \pm \gamma_5)b|B^-\rangle = i\bar{u}(p_{\Lambda})[(F_A^{\Lambda \bar{p}} \not{p}\gamma_5 \pm F_V^{\Lambda \bar{p}} \not{p}) + (F_P^{\Lambda \bar{p}} \gamma_5 \pm F_S^{\Lambda \bar{p}})]v(p_{\bar{p}}),
\]

where the form factors related to those of \( B^- \to p\bar{p} \) in Eq. (7) are shown as

\[
F_A^{\Lambda \bar{p}} = \sqrt{\frac{3}{2}} \frac{3}{10} (F_V - F_A), \quad F_V^{\Lambda \bar{p}} = -\sqrt{\frac{3}{2}} \frac{3}{10} (F_V - F_A), \quad F_P^{\Lambda \bar{p}} = \sqrt{\frac{3}{2}} \frac{3}{4} F_P.
\]

The three form factors \( F_A, F_V \) and \( F_P \) can be simply presented as [19, 20]

\[
F_{A,V} = \frac{C_{A,V}}{t^3}, \quad F_P = \frac{C_P}{t^4},
\]

where \( C_i \) \((i = A,V,P)\) are new parametrized form factors, which are taking to be real.

From the relation \( p^\mu \langle \Lambda \bar{p}|\bar{s}\gamma_{\mu}(1 - \gamma_5)b|B^-\rangle = m_b \langle \Lambda \bar{p}|\bar{s}(1 - \gamma_5)b|B^-\rangle \) in the heavy b quark limit, the parameters \( a_i(c_i) \) in Eq. (5) are associated with the scalar and pseudo-scalar matrix elements defined in Eq. (5). As a result, we get that

\[
a_1 = m_b F_A^{\Lambda \bar{p}}, \quad a_3 = \frac{m_b F_P^{\Lambda \bar{p}}}{p \cdot (p_{\bar{p}} - p_{\Lambda})}, \quad c_1 = m_b F_V^{\Lambda \bar{p}}, \quad c_3 = \frac{m_b F_P^{\Lambda \bar{p}}}{p \cdot (p_{\bar{p}} + p_{\Lambda})}.
\]

\[4\]
The amplitude in Eq. (5) then becomes

\[
A(B^- \to \Lambda \bar{p} \gamma) = \frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* \frac{e}{8\pi^2} 2c_{eff}^2 \left\{ m_b^3 \varepsilon^\mu \bar{u}(p_A) \left[ F_A^{\Lambda \bar{p}} \gamma_\mu \gamma_5 + F_P^{\Lambda \bar{p}} \gamma_\mu \frac{(p_\bar{p} - p_A) \mu}{p \cdot (p_\bar{p} - p_A)} - F_V^{\Lambda \bar{p}} \gamma_\mu - F_P^{\Lambda \bar{p}} \frac{(p_\bar{p} + p_A) \mu}{p \cdot (p_\bar{p} + p_A)} \right] v(p_\bar{p}) \right. \\
-2m_b p_B \cdot \varepsilon \ u(p_A) \left[ F_A^{\Lambda \bar{p}} p_\gamma \gamma_5 + F_P^{\Lambda \bar{p}} p_\gamma + F_V^{\Lambda \bar{p}} \right] v(p_\bar{p}) \right\},
\]

(12)

with three unknown form factors \( F_A^{\Lambda \bar{p}} \), \( F_V^{\Lambda \bar{p}} \) and \( F_P^{\Lambda \bar{p}} \). We note that the terms corresponding to \( a_2 \) disappear due to the fact of \( \varepsilon \cdot p = 0 \). Even though \( c_2 \) can only be determined by experimental data, according to QCD counting rules, \( c_2 \) needs an additional \( 1/t \) than \( c_1 \) to flip the helicity, so that it is guaranteed to give a small contribution and can be neglected.

After summing over the photon polarizations and baryon spins, from Eq. (12), the decay rate of \( \Gamma \) is given by the integration of

\[
d\Gamma = \frac{1}{(2\pi)^3} \frac{m_b^6}{4M_B^2 E_\gamma^2} |C_t|^2 \left[ V|F_V^{\Lambda \bar{p}}|^2 + A|F_A^{\Lambda \bar{p}}|^2 + P|F_P^{\Lambda \bar{p}}|^2 + I_{VP} Re(F_V^{\Lambda \bar{p}} F_P^{\Lambda \bar{p}}^\ast) \right] dm_{\Lambda \bar{p}} dm_{\gamma}^2
\]

(13)

where

\[
m_{\Lambda \bar{p}} = p_A + p_\bar{p}, \quad m_{\gamma} = p_\gamma, \quad C_t = \frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* \frac{e}{8\pi^2} 2c_{eff}^2,
\]

\[
V(A) = p_A \cdot p(E_\gamma E_\gamma - p_\gamma \cdot p) + E_\gamma (E_A p_\bar{p} \cdot p \pm E_\gamma m_\Lambda m_\bar{p}),
\]

\[
P = -\frac{E_\gamma (E_A + E_\bar{p})(m_\Lambda m_\bar{p} - p_A \cdot p_\bar{p})}{p_A \cdot p + p_\bar{p} \cdot p} + \frac{(m_A^2 + m_\bar{p}^2 + 2p_A \cdot p_\bar{p})(m_\Lambda m_\bar{p} - p_A \cdot p_\bar{p})}{2(p_A \cdot p + p_\bar{p} \cdot p)^2}
\]+ \frac{E_\gamma (E_A - E_\bar{p})(m_\Lambda m_\bar{p} - p_A \cdot p_\bar{p})}{p_A \cdot p - p_\bar{p} \cdot p} - \frac{(m_A^2 + m_\bar{p}^2 - 2p_A \cdot p_\bar{p})(m_\Lambda m_\bar{p} + p_A \cdot p_\bar{p})}{2(p_A \cdot p - p_\bar{p} \cdot p)^2}
\]-p_A \cdot p_\bar{p},
\]

\[
I_{VP(AP)} = 2E_\bar{p} E_\gamma m_\Lambda - p_\bar{p} \cdot p m_\Lambda \pm E_A E_\gamma (m_\Lambda - m_\bar{p}) \pm m_\Lambda p_A \cdot p
\]+ \frac{E_\gamma (E_\bar{p} \pm E_A)(m_\Lambda + m_\bar{p})p_A \cdot p - E_\gamma^2 (m_\Lambda - m_\bar{p})(p_A \cdot p_\bar{p} \pm m_\Lambda m_\bar{p})}{p_A \cdot p \pm p_\bar{p} \cdot p},
\]

(14)

It is important to note that, since the penguin-induced radiative \( B \) decays are associated
with axial-vector currents shown in Eq. (5), we have used 21
\[
\sum_{\lambda=1,2} \varepsilon^{*\lambda}_{\mu} \varepsilon^{\lambda}_{\nu} = -g_{\mu\nu} + \frac{k_{\mu} n_{\nu} + k_{\nu} n_{\mu}}{k \cdot n} - \frac{k_{\mu} k_{\nu}}{(k \cdot n)^2},
\]
(15)
where \( n = (1, 0, 0, 0) \), to sum over the photon polarizations instead of the direct replacement of \( \sum_{\lambda=1,2} \varepsilon^{*\lambda}_{\mu} \varepsilon^{\lambda}_{\nu} \rightarrow -g_{\mu\nu} \) which is valid in the QED-like theory due to the Ward identity.

For the numerical analysis of the branching ratios, we take the effective Wilson coefficient \( c_7^{\text{eff}} = -0.314 \) 22, the running quark mass \( m_b = 4.88 \) GeV and CKM matrix elements \( V_{tb} V_{ts}^* = -0.0402 \). Even though there are no theoretical calculations to the unknown \( C_A, C_V \) and \( C_P \). By virtue of the approach of Ref. 20, these form factors are related to the present experimental data, such as \( Br(B^- \rightarrow p\bar{p}\pi^-) \), \( Br(B^0 \rightarrow p\bar{p}K^0) \), \( Br(B^- \rightarrow p\bar{p}K^-) \) 15 and \( Br(B^- \rightarrow \Lambda \bar{\Lambda}K^-) \) 23, characterized by an emitted pseudoscalar meson. For a reliable \( \chi^2 \) fitting, we need 2 degrees of freedom (DOF) by ignoring the \( C_P \) term since its contribution is always associated with one more \( 1/t \) over \( C_A \) and \( C_V \) ones, as seen in Eq. (10). We will take a consistent check in the next paragraph to this simplification. To illustrate our results, we fix the color number \( N_C = 3 \) and weak phase \( \gamma = 54.8^\circ \). The input experimental data and numerical values are summarized in Table I.

| Input | experimental data | Fit result | best fit (with 1 \( \sigma \) error) |
|-------|------------------|------------|-----------------------------------|
| \( Br(B^- \rightarrow p\bar{p}\pi^-) \) [15] | 3.06 \( \pm \) 0.82 | \( C_A \) | \(-68.3 \pm 5.1\) |
| \( Br(B^0 \rightarrow p\bar{p}K^0) \) [15] | 1.88 \( \pm \) 0.80 | \( C_V \) | 35.1 \( \pm \) 9.0 |
| \( Br(B^- \rightarrow p\bar{p}K^-) \) [15] | 5.66 \( \pm \) 0.91 | \( \chi^2/DOF \) | 1.85 |
| \( Br(B^- \rightarrow \Lambda \bar{\Lambda}K^-) \) [23] | 2.91 \( \pm \) 0.98 | | |

Using the fitted values of \( C_A \) and \( C_V \), we find \( Br(B^- \rightarrow \Lambda \bar{p}\gamma) = (0.92 \pm 0.20) \times 10^{-6} \) which is larger than its two-body decay partner as expected and it is close to the result of \( 1.2 \times 10^{-6} \) in the pole model 3. However, our predicted value on \( B^- \rightarrow \Lambda \bar{p}\gamma \) is smaller than \( (2.16^{+0.58}_{-0.33} \pm 0.20) \times 10^{-6} \) 24 measured by BELLE. If we put this new observed value into our fitting, we can further include \( C_P \) ignored previously. The fitted values are \( C_A = -73.3 \pm 9.1 \) GeV\(^4\), \( C_V = 43.7 \pm 12.1 \) GeV\(^4\) and \( C_P = 134.3 \pm 327.0 \) GeV\(^7\) with \( \chi^2/DOF = 3.65 \) which is about two times bigger than previous one. Clearly, it presents an inferior fitting with small \( C_{A,V} \) changes. When putting back these three fitted values to \( Br(B^- \rightarrow \Lambda \bar{p}\gamma) \) for
FIG. 2: $dBr(B^{-} \rightarrow \Lambda\bar{p}\gamma)/dm_{\Lambda\bar{p}}$ vs. $m_{\Lambda\bar{p}}$. The solid line stands for the input values of $(C_{A}, C_{V}) = (-68.3, 35.1)$ while the dash line stands for those of $(C_{A}, C_{V}, C_{P}) = (-73.3, 43.7, 134.3)$. A consistency check, we get $(1.16 \pm 0.31) \times 10^{-6}$ regardless of inputting larger experimental value, which explains the large value of $\chi^{2}/DOF$. The insensitivity of $C_{P}$ on the decay branching ratio justifies our early simplification of ignoring its contribution beside the $1/t$ argument.

In Ref. [1], it was suggested that the reduced energy release can make the branching ratios of three-body decays as significant as their counterparts of two-body modes or even larger, and one of the signatures would be baryon pair threshold effect [1, 20]. In Fig. 2 from Eq. (13) we show the differential branching ratio of $dBr(B^{-} \rightarrow \Lambda\bar{p}\gamma)/dm_{\Lambda\bar{p}}$ representing the threshold enhancement around the invariant mass $m_{\Lambda\bar{p}} = 2.05$ GeV, which is consistent with Fig. 2 in Ref. [24] of the BELLE result. Around the threshold, the baryon pair contains half of the $B$ meson energy while the phone emitting back to back to the baryon pair with another half of energy which explains the peak at $E_{\gamma} \sim 2$ GeV in Fig. 3 of Ref. [24]. Such mechanism is similar to the two-body decays so that factorization method works [1] even in the three-body decays.

To discuss other radiative baryonic $B^{-}$ decays, we give form factors by relating them to $F_{V,A,P}$ in the $B^{-} \rightarrow p\bar{p}$ transition similar to the case of $B^{-} \rightarrow \Lambda\bar{p}\gamma$ as follows:

\[ B^{-} \rightarrow \Sigma^{0}\bar{p}\gamma : \]

\[ F_{V}^{\Sigma^{0}\bar{p}} = -\frac{11F_{V}}{10\sqrt{2}} - \frac{9F_{A}}{10\sqrt{2}} \quad F_{A}^{\Sigma^{0}\bar{p}} = -\frac{9F_{V}}{10\sqrt{2}} - \frac{11F_{A}}{10\sqrt{2}} \quad F_{P}^{\Sigma^{0}\bar{p}} = \frac{F_{P}}{3\sqrt{2}}, \]
To calculate the branching ratio of $B^- \rightarrow \Sigma^- \bar{n} \gamma$, we can use the formula in Eq. (13) by replacing $\Lambda$ and $\bar{p}$ by $B$ and $\bar{B}'$, respectively. The two sets of predicted values for $B^- \rightarrow B \bar{B}' \gamma$ with and without $C_P$ are shown in Table II respectively. As a comparison, we also list the work of the pole model approach by Cheng and Yang [9] in the table. We note that, in

\begin{equation}
\begin{aligned}
B^- \rightarrow \Sigma^- \bar{n} \gamma: \\
F_V^{\Sigma^- \bar{n}} = -\frac{11F_V}{10} - \frac{9F_A}{10}, \quad F_A^{\Sigma^- \bar{n}} = -\frac{9F_V}{10} - \frac{9F_A}{11}, \quad F_P^{\Sigma^- \bar{n}} = \frac{F_P}{4},
\end{aligned}
\end{equation}

\begin{equation}
\begin{aligned}
B^- \rightarrow \Xi^- \bar{\Lambda} \gamma: \\
F_V^{\Xi^- \bar{\Lambda}} = -\frac{21F_V}{10\sqrt{6}} - \frac{9F_A}{10\sqrt{6}}, \quad F_A^{\Xi^- \bar{\Lambda}} = -\frac{9F_V}{10\sqrt{6}} - \frac{21F_A}{10\sqrt{6}}, \quad F_P^{\Xi^- \bar{\Lambda}} = \frac{F_P}{4},
\end{aligned}
\end{equation}

\begin{equation}
\begin{aligned}
B^- \rightarrow \Xi^0 \bar{\Sigma}^- \gamma: \\
F_V^{\Xi^0 \bar{\Sigma}^-} = -\frac{F_V}{10}, \quad F_A^{\Xi^0 \bar{\Sigma}^-} = -\frac{9F_V}{10}, \quad F_P^{\Xi^0 \bar{\Sigma}^-} = \frac{5F_P}{4},
\end{aligned}
\end{equation}

\begin{equation}
\begin{aligned}
B^- \rightarrow \Xi^- \bar{\Sigma}^0 \gamma: \\
F_V^{\Xi^- \bar{\Sigma}^0} = -\frac{F_V}{10\sqrt{2}}, \quad F_A^{\Xi^- \bar{\Sigma}^0} = -\frac{9F_V}{10\sqrt{2}}, \quad F_P^{\Xi^- \bar{\Sigma}^0} = \frac{5F_P}{4\sqrt{2}}.
\end{aligned}
\end{equation}

To calculate the branching ratio of $B^- \rightarrow B \bar{B}' \gamma$, we can use the formula in Eq. (13) by replacing $\Lambda$ and $\bar{p}$ by $B$ and $\bar{B}'$, respectively. The two sets of predicted values for $B^- \rightarrow B \bar{B}' \gamma$ with and without $C_P$ are shown in Table II respectively. As a comparison, we also list the work of the pole model approach by Cheng and Yang [9] in the table. We note that, in

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|}
\hline
Branching Ratios & Fits & Pole model [9] \\
\hline
& $(C_A, C_V) = (-68.3 \pm 5.1, 35.0 \pm 9.0)$ & $(C_A, C_V, C_P) = (-73.3 \pm 9.1, 43.7 \pm 12.1, 134.3 \pm 327.0)$ \\
& $(0.92 \pm 0.20) \times 10^{-6}$ & $[(1.16 \pm 0.31) \times 10^{-6}]$ \\
$Br(B^- \rightarrow \Lambda \bar{p} \gamma)$ & $[(1.16 \pm 0.31) \times 10^{-6}]$ & $1.2 \times 10^{-6}$ \\
& $(1.7 \pm 1.5) \times 10^{-7}$ & $(1.2 \pm 1.2) \times 10^{-7}$ \\
$Br(B^- \rightarrow \Sigma^0 \bar{p} \gamma)$ & $(1.2 \pm 1.2) \times 10^{-7}$ & $2.9 \times 10^{-9}$ \\
& $(3.4 \pm 2.8) \times 10^{-7}$ & $(2.5 \pm 2.4) \times 10^{-7}$ \\
$Br(B^- \rightarrow \Sigma^- \bar{n} \gamma)$ & $(2.5 \pm 2.4) \times 10^{-7}$ & $5.7 \times 10^{-9}$ \\
& $(0.48 \pm 0.50) \times 10^{-7}$ & $(0.61 \pm 0.60) \times 10^{-7}$ \\
$Br(B^- \rightarrow \Xi^- \bar{\Lambda} \gamma)$ & $(0.61 \pm 0.60) \times 10^{-7}$ & $2.4 \times 10^{-7}$ \\
& $(3.3 \pm 0.7) \times 10^{-7}$ & $(3.7 \pm 0.9) \times 10^{-7}$ \\
$Br(B^- \rightarrow \Xi^0 \bar{\Sigma}^- \gamma)$ & $(3.7 \pm 0.9) \times 10^{-7}$ & $1.2 \times 10^{-6}$ \\
& $(1.5 \pm 0.6) \times 10^{-7}$ & $(1.8 \pm 0.6) \times 10^{-7}$ \\
$Br(B^- \rightarrow \Xi^- \bar{\Sigma}^0 \gamma)$ & $(1.8 \pm 0.6) \times 10^{-7}$ & $6.0 \times 10^{-7}$ \\
\hline
\end{tabular}
\caption{Decay branching ratios}
\end{table}

Table II the value in the bracket of the third column for $Br(B^- \rightarrow \Lambda \bar{p} \gamma)$ is not a prediction but a consistency comparison with the putting-back form factors, since we have used the observed value of $Br(B^- \rightarrow \Lambda \bar{p} \gamma)$ from BELLE. We found that, except for $Br(B^- \rightarrow \Lambda \bar{p} \gamma)$,
all predicted values are $O(10^{-7})$. In terms of inverse sign between $C_A$ and $C_V$, there are constructive effects for $F_A^{A0}$ and $F_V^{A0}$, which are proportional to $(F_V - F_A)$ as shown in Eq. (11), whereas destructive effects make other $F_A^{BB'}$ and $F_V^{BB'}$ in Eq. (10) small. Consequently, all modes for $B^-$ radiative baryonic decays are suppressed except for $Br(B^- \to \Lambda \bar{p} \gamma)$. We remark that such suppressions exist only in the SM-like theories. Thus, these radiative baryonic decays are useful modes for testing the new physics.

As seen in Table 11, both our results and those of the pole model satisfy the relations of $Br(B^- \to \Sigma^- \bar{n} \gamma) \simeq 2Br(B^- \to \Sigma^0 \bar{p} \gamma)$ and $Br(B^- \to \Xi^0 \Sigma^- \gamma) \simeq 2Br(B^- \to \Xi^- \Sigma^0 \gamma)$ because of the SU(3) symmetry. In the pole model, the decay branching ratios of $B^- \to \Lambda \bar{p} \gamma$ and $B^- \to \Xi^0 \Sigma^- \gamma$ are found to be large, around $1.2 \times 10^{-6}$, since there are intermediated elements of the $B \to p \gamma$ transition. Most of the predicted values for $Br(B \to BB')$ are spanning in the order of $10^{-7}$, which are larger than the estimated values of $O(10^{-9})$ due to the IB effects of their two-body counterparts. In particular, we have found that
$Br(B^- \rightarrow \Lambda\bar{p}\gamma)$ is $(1.16 \pm 0.31) \times 10^{-6}$ and $(0.92 \pm 0.20) \times 10^{-6}$ with and without $C_P$, respectively, which are consistent with the pole model prediction \cite{9} but smaller than the experimental data from BELLE \cite{24}. More precise measurements are clearly needed.

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