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Radiative baryonic B decays

C.Q. Geng, Y.K. Hsiao

Department of Physics, National Tsing Hua University, Hsinchu 300, Taiwan

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Abstract

We study the structure-dependent contributions to the radiative baryonic B decays of $B \to \boldsymbol{B}\bar{\boldsymbol{B}}'\gamma$ in the standard model. We show that the decay branching ratios of $Br(B \to \boldsymbol{B}\bar{\boldsymbol{B}}'\gamma)$ 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 1×10^{-6}, which is close to the pole model estimation but smaller than the experimental measurement from Belle.

The radiative baryonic B decays of $B \to \boldsymbol{B}\bar{\boldsymbol{B}}'\gamma$ are of interest since they are three-body decays with two spin-1/2 baryons ($\boldsymbol{B}$ and $\boldsymbol{B}'$) 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 $B \to \boldsymbol{B}\bar{\boldsymbol{B}}'\gamma$ are from their two-body decay counterparts of $B \to \boldsymbol{B}\bar{\boldsymbol{B}}'$ 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 $\alpha_{\rm em}$ comparing with their counterparts. According to the existing upper bounds of $B \to \boldsymbol{B}\bar{\boldsymbol{B}}'$, given by [6–8]

$$Br(\bar{B}^0 \to p\bar{p}) < 2.7 \times 10^{-7} \text{ (BaBar)}, \qquad Br(\bar{B}^0 \to \Lambda\bar{\Lambda}) < 7.9 \times 10^{-7} \text{ (Belle)},$$
$$Br(B^- \to \Lambda\bar{p}) < 4.6 \times 10^{-7} \text{ (Belle)}, \tag{1}$$

E-mail address: ykhsiao@phys.nthu.edu.tw (Y.K. Hsiao).

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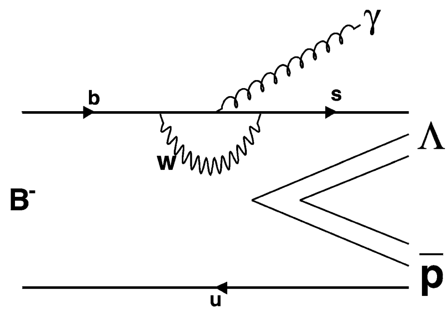

Fig. 1. Diagram for $B^- \to \Lambda\bar{p}\gamma$.

one finds that

$$Br(B \to \boldsymbol{B}\bar{\boldsymbol{B}}'\gamma)_{\mathrm{IB}} \leqslant O(10^{-9}). \tag{2}$$

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 $Br(B \to \boldsymbol{B}\bar{\boldsymbol{B}}'\gamma)$, 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 $Br(B^- \to \boldsymbol{B}\bar{\boldsymbol{B}}'\gamma)$ could be as large as $Br(B^- \to \boldsymbol{B}\bar{\boldsymbol{B}}')$. In this Letter, we shall concentrate on the SD contributions to $Br(B \to \boldsymbol{B}\bar{\boldsymbol{B}}'\gamma)$.

To start our study, we must tackle the cumbersome transition matrix elements in $B \to \boldsymbol{B}\bar{\boldsymbol{B}}'$. As more and more experimental data on three-body decays [13–15] 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 [18–20] 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 Letter, we handle the transition matrix elements according to the QCD counting rules.

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

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

with the tensor operator

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

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

$$A(B^- \to \Lambda\bar{p}\gamma) = \frac{G_F}{\sqrt{2}} V_{ts}^* V_{tb} \frac{e}{8\pi^2} 2c_7^{\mathrm{eff}} \bigl\{ m_b^2 \varepsilon^\mu \langle \Lambda\bar{p}|\bar{s}\gamma_\mu(1-\gamma_5)b|B^-\rangle - 2m_b p_B \cdot \varepsilon \langle \Lambda\bar{p}|\bar{s}(1+\gamma_5)b|B^-\rangle \bigr\}, \tag{5}$$

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

$$\begin{aligned} \langle \Lambda\bar{p}|\bar{s}\gamma_\mu b|B^-\rangle &= i\bar{u}(p_\Lambda)\bigl[a_1\gamma_\mu\gamma_5 + a_2 p_\mu\gamma_5 + a_3(p_{\bar{p}} - p_\Lambda)_\mu\gamma_5\bigr]v(p_{\bar{p}}), \\ \langle \Lambda\bar{p}|\bar{s}\gamma_\mu\gamma_5 b|B^-\rangle &= i\bar{u}(p_\Lambda)\bigl[c_1\gamma_\mu + c_2 i\sigma_{\mu\nu}p^\nu + c_3(p_{\bar{p}} + p_\Lambda)_\mu\bigr]v(p_{\bar{p}}), \end{aligned} \tag{6}$$

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. (6), 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 $B \to \boldsymbol{B}\bar{\boldsymbol{B}}'$ transitions based on the QCD counting rules [18], that require the form factors to behave as inverse powers of $t = (p_{\boldsymbol{B}} + p_{\bar{\boldsymbol{B}}'})^2$. The detail discussions can be referred to Refs. [19,20]. In this Letter, we shall follow their approach. The representations of the matrix elements for the $B^- \to p\bar{p}$ transition are given by [20]

$$\langle p\bar{p}|\bar{u}(1\pm\gamma_5)b|B^-\rangle = i\bar{u}(p_p)\big[(F_A \not{p}\gamma_5 \pm F_V \not{p}) + (F_P\gamma_5 \pm F_S)\big]v(p_{\bar{p}}), \tag{7}$$

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

$$\langle \Lambda\bar{p}|\bar{s}(1\pm\gamma_5)b|B^-\rangle = i\bar{u}(p_\Lambda)\big[\big(F_A^{\Lambda\bar{p}} \not{p}\gamma_5 \pm F_V^{\Lambda\bar{p}} \not{p}\big) + \big(F_P^{\Lambda\bar{p}}\gamma_5 \pm F_S^{\Lambda\bar{p}}\big)\big]v(p_{\bar{p}}), \tag{8}$$

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

$$F_A^{\Lambda\bar{p}} = \sqrt{\frac{3}{2}}\frac{3}{10}(F_V - F_A), \qquad F_V^{\Lambda\bar{p}} = -\sqrt{\frac{3}{2}}\frac{3}{10}(F_V - F_A), \qquad F_{P(S)}^{\Lambda\bar{p}} = \sqrt{\frac{3}{2}}\frac{3}{4}F_P. \tag{9}$$

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

$$F_{A,V} = \frac{C_{A,V}}{t^3}, \qquad F_P = \frac{C_P}{t^4}, \tag{10}$$

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. (6) are associated with the scalar and pseudo-scalar matrix elements defined in Eq. (8). As a result, we get that

$$a_1 = m_b F_A^{\Lambda\bar{p}}, \qquad a_3 = \frac{m_b F_P^{\Lambda\bar{p}}}{p\cdot(p_{\bar{p}} - p_\Lambda)}, \qquad c_1 = m_b F_V^{\Lambda\bar{p}}, \qquad c_3 = \frac{m_b F_P^{\Lambda\bar{p}}}{p\cdot(p_{\bar{p}} + p_\Lambda)}. \tag{11}$$

The amplitude in Eq. (5) then becomes

$$\begin{aligned} A(B^- \to \Lambda\bar{p}\gamma) = \frac{G_F}{\sqrt{2}}V_{tb}V_{ts}^*\frac{e}{8\pi^2}2c_7^{\mathrm{eff}}\bigg\{ & m_b^3\varepsilon^\mu\bar{u}(p_\Lambda)\bigg[F_A^{\Lambda\bar{p}}\gamma_\mu\gamma_5 + F_P^{\Lambda\bar{p}}\gamma_5\frac{(p_{\bar{p}} - p_\Lambda)_\mu}{p\cdot(p_{\bar{p}} - p_\Lambda)} - F_V^{\Lambda\bar{p}}\gamma_\mu \\ & - F_P^{\Lambda\bar{p}}\frac{(p_{\bar{p}} + p_\Lambda)_\mu}{p\cdot(p_{\bar{p}} + p_\Lambda)}\bigg]v(p_{\bar{p}}) \\ & - 2m_b p_B\cdot\varepsilon u(p_\Lambda)\Big[F_A^{\Lambda\bar{p}}\not{p}\gamma_5 + F_P^{\Lambda\bar{p}}\gamma_5 + F_V^{\Lambda\bar{p}}\not{p} + F_P^{\Lambda\bar{p}}\Big]v(p_{\bar{p}})\bigg\}, \end{aligned} \tag{12}$$

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 Γ is given by the integration of

$$\begin{aligned} d\Gamma = {} & \frac{1}{(2\pi)^3}\frac{m_b^6}{4M_B^3E_\gamma^2}|C_t|^2 \\ & \times\big[V\big|F_V^{\Lambda\bar{p}}\big|^2 + A\big|F_A^{\Lambda\bar{p}}\big|^2 + P\big|F_P^{\Lambda\bar{p}}\big|^2 + I_{VP}\,\mathrm{Re}\big(F_V^{\Lambda\bar{p}}F_P^{\Lambda\bar{p}*}\big) + I_{AP}\,\mathrm{Re}\big(F_A^{\Lambda\bar{p}}F_P^{\Lambda\bar{p}*}\big)\big]\,dm_{\Lambda\bar{p}}^2\,dm_{\bar{p}\gamma}^2, \end{aligned} \tag{13}$$

where

$$
\begin{aligned}
& m_{\Lambda\bar{p}} = p_\Lambda + p_{\bar{p}}, \qquad m_{\bar{p}\gamma} = p_{\bar{p}} + p_\gamma, \qquad C_t = \frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* \frac{e}{8\pi^2} 2c_7^{\text{eff}}, \\
& V(A) = p_\Lambda \cdot p(E_{\bar{p}} E_\gamma - p_{\bar{p}} \cdot p) + E_\gamma (E_\Lambda p_{\bar{p}} \cdot p \pm E_\gamma m_\Lambda m_{\bar{p}}), \\
& P = -\frac{E_\gamma (E_\Lambda + E_{\bar{p}})(m_\Lambda m_{\bar{p}} - p_\Lambda \cdot p_{\bar{p}})}{p_\Lambda \cdot p + p_{\bar{p}} \cdot p} + \frac{(m_\Lambda^2 + m_{\bar{p}}^2 + 2 p_\Lambda \cdot p_{\bar{p}})(m_\Lambda m_{\bar{p}} - p_\Lambda \cdot p_{\bar{p}})}{2(p_\Lambda \cdot p + p_{\bar{p}} \cdot p)^2} \\
& \quad + \frac{E_\gamma (E_\Lambda - E_{\bar{p}})(m_\Lambda m_{\bar{p}} - p_\Lambda \cdot p_{\bar{p}})}{p_\Lambda \cdot p - p_{\bar{p}} \cdot p} - \frac{(m_\Lambda^2 + m_{\bar{p}}^2 - 2 p_\Lambda \cdot p_{\bar{p}})(m_\Lambda m_{\bar{p}} + p_\Lambda \cdot p_{\bar{p}})}{2(p_\Lambda \cdot p - p_{\bar{p}} \cdot p)^2} - p_\Lambda \cdot p_{\bar{p}}, \\
& I_{VP(AP)} = 2 E_{\bar{p}} E_\gamma m_\Lambda - p_{\bar{p}} \cdot p m_\Lambda \pm E_\Lambda E_\gamma (m_\Lambda - m_{\bar{p}}) \pm m_{\bar{p}} p_\Lambda \cdot p \\
& \qquad + \frac{E_\gamma (E_{\bar{p}} \pm E_\Lambda)(m_\Lambda + m_{\bar{p}}) p_\Lambda \cdot p - E_\gamma^2 (m_\Lambda - m_{\bar{p}})(p_\Lambda \cdot p_{\bar{p}} \pm m_\Lambda m_{\bar{p}})}{p_\Lambda \cdot p \pm p_{\bar{p}} \cdot p}.
\end{aligned}
\tag{14}
$$

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]

$$
\sum_{\lambda=1,2} \varepsilon_\mu^{*\lambda} \varepsilon_\nu^\lambda = -g_{\mu\nu} + \frac{k_\mu n_\nu + k_\nu n_\mu}{k \cdot n} - \frac{k_\mu k_\nu}{(k \cdot n)^2}, \tag{15}
$$

where $n = (1, 0, 0, 0)$, to sum over the photon polarizations instead of the direct replacement of $\sum_{\lambda=1,2} \varepsilon_\mu^{*\lambda} \varepsilon_\nu^\lambda \to -g_{\mu\nu}$ 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 $Br(B^- \to p\bar{p}\pi^-)$, $Br(B^0 \to p\bar{p}K^0)$, $Br(B^- \to p\bar{p}K^-)$ [15] and $Br(B^- \to \Lambda\bar{\Lambda}K^-)$ [23], characterized by an emitted pseudoscalar meson. For a reliable χ^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 $\gamma = 54.8°$. The input experimental data and numerical values are summarized in Table 1.

Using the fitted values of C_A and C_V, we find $Br(B^- \to \Lambda\bar{p}\gamma) = (0.92 \pm 0.20) \times 10^{-6}$ which is larger than its two-body decay partner as expected and it is close to the result of 1.2×10^{-6} in the pole model [9]. However, our predicted value on $B^- \to \Lambda\bar{p}\gamma$ is smaller than $(2.16^{+0.58}_{-0.53} \pm 0.20) \times 10^{-6}$ [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 $C_A = -73.3 \pm 9.1\ \text{GeV}^4$, $C_V = 43.7 \pm 12.1\ \text{GeV}^4$ and $C_P = 134.3 \pm 327.0\ \text{GeV}^7$ with $\chi^2/\text{DOF} = 3.65$ 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 $Br(B^- \to \Lambda\bar{p}\gamma)$ for 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 χ^2/DOF. The insensitivity

Table 1
Fits of C_A, C_V in units of GeV^4

Input	Experimental data	Fit result	Best fit (with 1σ error)
$Br(B^- \to p\bar{p}\pi^-)$ [15]	3.06 ± 0.82	C_A	-68.3 ± 5.1
$Br(B^0 \to p\bar{p}K^0)$ [15]	1.88 ± 0.80	C_V	35.1 ± 9.0
$Br(B^- \to p\bar{p}K^-)$ [15]	5.66 ± 0.91	χ^2/DOF	1.85
$Br(B^- \to \Lambda\bar{\Lambda}K^-)$ [23]	2.91 ± 0.98		

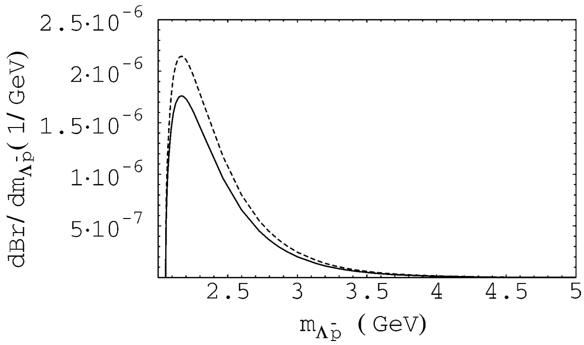

Fig. 2. $dBr(B^- \to \Lambda \bar{p}\gamma)/dm_{\Lambda\bar{p}}$ vs. $m_{\Lambda\bar{p}}$. The solid line stands for the input values of $(C_A, C_V) = (-68.3, 35.1)$ while the dash line stands for those of $(C_A, C_V, C_P) = (-73.3, 43.7, 134.3)$.

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 $dBr(B^- \to \Lambda\bar{p}\gamma)/dm_{\Lambda\bar{p}}$ vs. $m_{\Lambda\bar{p}}$ 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 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 $B^- \to p\bar{p}$ transition similar to the case of $B^- \to \Lambda\bar{p}\gamma$ as follows:

$$
\begin{array}{llll}
B^- \to \Sigma^0\bar{p}\gamma: & F_V^{\Sigma^0\bar{p}} = -\dfrac{11F_V}{10\sqrt{2}} - \dfrac{9F_A}{10\sqrt{2}}, & F_A^{\Sigma^0\bar{p}} = -\dfrac{9F_V}{10\sqrt{2}} - \dfrac{11F_A}{10\sqrt{2}}, & F_P^{\Sigma^0\bar{p}} = \dfrac{F_P}{3\sqrt{2}}, \\
B^- \to \Sigma^-\bar{n}\gamma: & F_V^{\Sigma^-\bar{n}} = -\dfrac{11F_V}{10} - \dfrac{9F_A}{10}, & F_A^{\Sigma^-\bar{n}} = -\dfrac{9F_V}{10} - \dfrac{9F_A}{11}, & F_P^{\Sigma^-\bar{n}} = \dfrac{F_P}{4}, \\
B^- \to \Xi^-\bar{\Lambda}\gamma: & F_V^{\Xi^-\bar{\Lambda}} = -\dfrac{21F_V}{10\sqrt{6}} - \dfrac{9F_A}{10\sqrt{6}}, & F_A^{\Xi^-\bar{\Lambda}} = -\dfrac{9F_V}{10\sqrt{6}} - \dfrac{21F_A}{10\sqrt{6}}, & F_P^{\Xi^-\bar{\Lambda}} = \dfrac{F_P}{4}, \\
B^- \to \Xi^0\bar{\Sigma}^-\gamma: & F_V^{\Xi^0\bar{\Sigma}^-} = -\dfrac{F_V}{10} - \dfrac{9F_A}{10}, & F_A^{\Xi^0\bar{\Sigma}^-} = -\dfrac{9F_V}{10} - \dfrac{F_A}{10}, & F_P^{\Xi^0\bar{\Sigma}^-} = \dfrac{5F_P}{4}, \\
B^- \to \Xi^-\bar{\Sigma}^0\gamma: & F_V^{\Xi^-\bar{\Sigma}^0} = -\dfrac{F_V}{10\sqrt{2}} - \dfrac{9F_A}{10\sqrt{2}}, & F_A^{\Xi^-\bar{\Sigma}^0} = -\dfrac{9F_V}{10\sqrt{2}} - \dfrac{F_A}{10\sqrt{2}}, & F_P^{\Xi^-\bar{\Sigma}^0} = \dfrac{5F_P}{4\sqrt{2}}.
\end{array}
\tag{16}
$$

To calculate the branching ratio of $B \to \boldsymbol{B}\bar{\boldsymbol{B}}'\gamma$, we can use the formula in Eq. (13) by replacing Λ and $\bar{p}$ by $\boldsymbol{B}$ and $\bar{\boldsymbol{B}}'$, respectively. The two sets of predicted values for $B \to \boldsymbol{B}\bar{\boldsymbol{B}}'\gamma$ with and without C_P are shown in Table 2, 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 Table 2, the value in the bracket of the third column for $Br(B^- \to \Lambda\bar{p}\gamma)$ is not a prediction but a consistency comparison with the putting-back form factors, since we have used the observed value of $Br(B^- \to \Lambda\bar{p}\gamma)$ from Belle. We found that, except for $Br(B^- \to \Lambda\bar{p}\gamma)$, all predicted values are $O(10^{-7})$. In terms of inverse sign between C_A and C_V, there are constructive effects for $F_A^{\Lambda\bar{p}}$ and $F_V^{\Lambda\bar{p}}$, which are proportional to $(F_V - F_A)$ as shown in Eq. (9), whereas destructive effects make other $F_A^{\boldsymbol{B}\bar{\boldsymbol{B}}'}$ and $F_V^{\boldsymbol{B}\bar{\boldsymbol{B}}'}$ in Eq. (16) small.

Table 2
Decay branching ratios

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

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 2, 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 $Br(B^- \to \Xi^0\bar{\Sigma}^-\gamma) \simeq 2Br(B^- \to \Xi^-\bar{\Sigma}^0\gamma)$ because of the SU(3) symmetry. In the pole model, the decay branching ratios of $B^- \to \Lambda\bar{p}\gamma$ and $B^- \to \Xi^0\bar{\Sigma}^-\gamma$ are found to be large, around 1.2×10^{-6}, since they are intermediated through Λ_b and Ξ_b, which correspond to large coupling constants $g_{\Lambda_b \to B^- p}$ and $g_{\Xi_b^0 \to B^- \Sigma^+}$, respectively. However, in our work, the branching ratio of $B^- \to \Lambda\bar{p}\gamma$ is about three times larger than that of $B^- \to \Xi^0\bar{\Sigma}^-\gamma$, which is $O(10^{-7})$. Regardless of these differences, both two methods are within the experimental data allowed ranges, such as those of

$$\left[Br(B^- \to \Lambda\bar{p}\gamma) + 0.3\,Br(B^- \to \Sigma^0\bar{p}\gamma)\right]_{E_\gamma > 2.0\ \mathrm{GeV}} < 3.3 \times 10^{-6},$$

$$\left[Br(B^- \to \Sigma^0\bar{p}\gamma) + 0.4\,Br(B^- \to \Lambda\bar{p}\gamma)\right]_{E_\gamma > 2.0\ \mathrm{GeV}} < 6.4 \times 10^{-6},$$

from CLEO [25] and $Br(B^- \to \Sigma^0\bar{p}\gamma) < 3.3 \times 10^{-6}$ from Belle [24].

Finally, we relate the $\bar{B}^0$ decays with the corresponding B^- modes in terms of QCD counting rules even though there are no experimental data on radiative baryonic $\bar{B}^0$ decays. When neglecting the mass and life time differences, we obtain

$$\begin{aligned} &Br(B^- \to \Lambda\bar{p}\gamma) = Br(\bar{B}^0 \to \Lambda\bar{n}\gamma), && Br(B^- \to \Sigma^0\bar{p}\gamma) = Br(\bar{B}^0 \to \Sigma^0\bar{n}\gamma),\\ &Br(B^- \to \Sigma^-\bar{n}\gamma) = Br(\bar{B}^0 \to \Sigma^+\bar{p}\gamma), && Br(B^- \to \Xi^-\bar{\Lambda}\gamma) = Br(\bar{B}^0 \to \Xi^0\bar{\Lambda}\gamma),\\ &Br(B^- \to \Xi^-\bar{\Sigma}^0\gamma) = Br(\bar{B}^0 \to \Xi^0\bar{\Sigma}^0\gamma), && Br(B^- \to \Xi^0\bar{\Sigma}^-\gamma) = Br(\bar{B}^0 \to \Xi^-\bar{\Sigma}^+\gamma), \end{aligned} \tag{17}$$

which are also guaranteed by the SU(3) symmetry. From Eq. (17), we see that $Br(\bar{B}^0 \to \Lambda\bar{n}\gamma)$ can be as large as $Br(B^- \to \Lambda\bar{p}\gamma)$.

In sum, we have shown that the SD contributions to the radiative baryonic decays of $B \to \boldsymbol{B}\bar{\boldsymbol{B}}'\gamma$ in the SM are associated with the form factors of F_A, F_V and F_P in the matrix elements of the $B^- \to p\bar{p}$ transition. Most of the predicted values for $Br(B \to \boldsymbol{B}\bar{\boldsymbol{B}}'\gamma)$ 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 $Br(B^- \to \Lambda\bar{p}\gamma)$ 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 [9] but smaller than the experimental data from Belle [24]. More precise measurements are clearly needed.

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