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Nonperturbative charming penguin contributions to isospin asymmetries in radiative B decays

Chul Kim,[1,*] Adam K. Leibovich,[2,+] and Thomas Mehen[1,‡]

[1]*Department of Physics, Duke University, Durham, North Carolina 27708, USA*

[2]*Department of Physics and Astronomy, University of Pittsburgh, Pennsylvania 15260, USA*

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Recent experimental data on the radiative decays $B \to V\gamma$, where V is a light vector meson, find small isospin violation in $B \to K^*\gamma$ while isospin asymmetries in $B \to \rho\gamma$ are of order 20%, with large uncertainties. Using soft-collinear effective theory, we calculate isospin asymmetries in these radiative B decays up to $O(1/m_b)$, also including $O(v\alpha_s)$ contributions from nonperturbative charming penguins (NPCP). In the absence of NPCP contributions, the theoretical predictions for the asymmetries are a few percent or less. Including the NPCP can significantly increase the isospin asymmetries for both $B \to V\gamma$ modes. We also consider the effect of the NPCP on the branching ratio and CP asymmetries in $B^\pm \to V^\pm\gamma$.

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The rare radiative B decays $B \to V\gamma$, where V is a light vector meson, are important in heavy flavor physics because the dominant processes are due to the flavor changing neutral current. Isospin asymmetries are interesting observables for testing the standard model (SM) and investigating new physics in the flavor sector. The isospin asymmetries for $B \to K^*\gamma$ and $B \to \rho\gamma$ are defined to be

$$\Delta_{0-}^{K^*} = \frac{\Gamma(\bar{B}^0 \to \bar{K}^{*0}\gamma) - \Gamma(B^- \to K^{*-}\gamma)}{\Gamma(\bar{B}^0 \to \bar{K}^{*0}\gamma) + \Gamma(B^- \to K^{*-}\gamma)},$$
$$\Delta_{0-}^{\rho} = \frac{2\Gamma(\bar{B}^0 \to \rho^0\gamma) - \Gamma(B^- \to \rho^-\gamma)}{2\Gamma(\bar{B}^0 \to \rho^0\gamma) + \Gamma(B^- \to \rho^-\gamma)}. \qquad (1)$$

In these asymmetries, the decay rates are averaged over charge conjugate modes. Recent experimental measurements find [1]

$$\Delta_{0-}^{K^*} = 0.03 \pm 0.04, \qquad \Delta_{0-}^{\rho} = 0.26 \pm 0.14, \qquad (2)$$

where the average values for the decay rates for $B \to K^*(\rho)\gamma$ are taken from the Heavy Flavor Averaging Group [2]. The isospin asymmetry for $B \to K^*\gamma$ is consistent with zero within an error of a few percent. The data suggest that the asymmetry in $B \to \rho\gamma$ is significantly larger, but because of large uncertainties it is not yet possible to draw a definitive conclusion. The work in this paper is motivated by the question of whether an anomalously large isospin asymmetry in $B \to \rho\gamma$ can be understood within the SM. In particular, we calculate subleading contributions to the leading QCD factorization theorems for $B \to V\gamma$ to see if they can explain the observed asymmetries.

In the heavy quark limit, the leading amplitudes are factorizable [3–7]. However, isospin-violating asymmetries come from $O(1/m_b)$ suppressed power corrections for which the factorization is necessarily more complicated. In this paper we calculate $O(1/m_b)$ corrections to the asymmetries using soft-collinear effective theory (SCET) [8], which provides a systematic power counting. In addition, possible endpoint divergences in these higher order corrections can be regulated without imposing an arbitrary infrared cutoff by including the zero-bin subtraction of Ref. [9]. Previous QCD analyses of isospin asymmetries in radiative B decays appear in Refs. [10–12]. The main difference between our analysis and previous work is the inclusion of nonperturbative charming penguin (NPCP) contributions, which are already known to play an important role in nonleptonic B decays [13–15]. (For an alternative point of view, see Ref. [16]).

The NPCP contributions to $B \to V\gamma$ are depicted in Fig. 1(a). In certain kinematic regimes, the invariant mass of the charm quark pair in the loop in Fig. 1(a) is near the threshold $2m_c$ in which case the charm quark pair is described by nonrelativistic QCD (NRQCD) and additional interactions need to be taken into account. As pointed out in Ref. [15], contributions from this regime are suppressed by only $v\alpha_s(2m_c)$ compared to the leading contribution. Here v is the relative velocity of the charm quarks in the threshold region. Therefore, the NPCP contribution to the isospin asymmetry could dominate over other $1/m_b$ suppressed contributions. In this paper, we

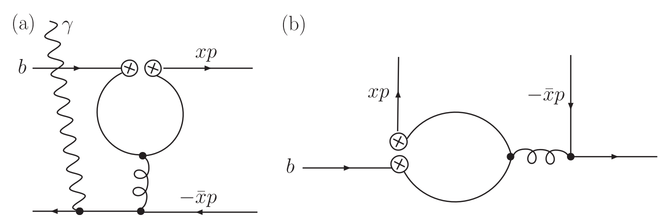

FIG. 1. Nonperturbative charming penguin (NPCP) contributions for (a) $B \to V\gamma$ and (b) $B \to M_1M_2$ arise when $\bar{x} = 1 - x \approx 4m_c^2/m_b^2$, in which case the charm quark pair is in the threshold region.

*Electronic address: chul@phy.duke.edu
+Electronic address: akl2@pitt.edu
‡Electronic address: mehen@phy.duke.edu

calculate the isospin asymmetries including the NPCP along with $1/m_b$ suppressed contributions. We also calculate the NPCP contributions to the branching ratio and CP asymmetries for $B^\pm \to V^\pm\gamma$.

In the absence of NPCP contributions, the theoretical predictions for $\Delta_{0-}^{K^*}$ and Δ_{0-}^{ρ} are no larger than a few percent. Including the NPCP contributions can significantly modify the theoretical predictions for the isospin asymmetries. We will see below that the NPCP contribution can be factorized using SCET, and the result expressed in terms of nonperturbative matrix elements. The NPCP matrix elements are fitted to available data on the isospin asymmetries, $\Delta_{0-}^{K^*}$ and Δ_{0-}^{ρ}, the CP asymmetry for $B^\pm \to \rho^\pm\gamma$ (recently measured by Belle [17]), and the branching ratio for $B^+ \to \rho^+\gamma$ [2]. The predictions for $\Delta_{0-}^{K^*}$ and Δ_{0-}^{ρ} are of order 10%, with uncertainties large enough that both predictions are consistent with experiment. However, the NPCP does not predict a large difference between $\Delta_{0-}^{K^*}$ and Δ_{0-}^{ρ}, as suggested by the central values in Eq. (2).

The isospin asymmetry in $B \to V\gamma$ can arise from either the mass difference of the spectator quark in the B-meson or the electric charge difference when the spectator quark emits the photon in the final state. However the isospin asymmetry due to the mass difference is negligible because it is $\mathcal{O}((m_u - m_d)/\Lambda_{\rm QCD})$ and therefore of order 1% or smaller. So the dominant piece comes from electromagnetic (EM) interactions with the spectator quark.

In order to describe the isospin-breaking corrections to $B \to V\gamma$, we need the following effective weak Hamiltonian

$$H_W = \frac{G_F}{\sqrt{2}}\bigg[\sum_{p=u,c}\lambda_p^{(q)}(C_1O_{1p} + C_2O_{2p}) - \lambda_t^{(q)}\bigg(\sum_{i=3}^{6}C_iO_i + C_{8g}O_{8g} + C_{7\gamma}O_{7\gamma}\bigg)\bigg], \quad (3)$$

where the operators are

$$\begin{aligned}
O_{1p} &= (\bar{p}b)_{V-A}(\bar{q}p)_{V-A},\\
O_{2p} &= (\bar{p}_\beta b_\alpha)_{V-A}(\bar{q}_\alpha p_\beta)_{V-A},\\
O_{3,5} &= (\bar{q}b)_{V-A}\sum_{q'=u,d,s,c,b}(\bar{q}'q')_{V\mp A},\\
O_{4,6} &= (\bar{q}_\beta b_\alpha)_{V-A}\sum_{q'=u,d,s,c,b}(\bar{q}'_\alpha q'_\beta)_{V\mp A},\\
O_{7\gamma} &= -\frac{em_b}{8\pi^2}\bar{q}\sigma^{\mu\nu}F_{\mu\nu}(1+\gamma_5)b,\\
O_{8g} &= -\frac{gm_b}{8\pi^2}\bar{q}\sigma^{\mu\nu}G^a_{\mu\nu}T^a(1+\gamma_5)b.
\end{aligned} \quad (4)$$

Here q is the d or s quark, the Cabibbo-Kobayashi-Maskawa (CKM) factor is $\lambda_p^{(q)} = V_{pb}V_{pq}^*$, and $V \pm A = \gamma^\mu(1 \pm \gamma_5)$.

For the asymmetric contributions, the photon radiates from either the initial or final spectator quark as shown in Fig. 2. If the photon is radiated from the initial spectator quark, we need SCET operators of the type $\mathcal{O}^{(0,4q)} = \bar{\chi}_{\bar{n}}\Gamma b_v\bar{\xi}_n\Gamma\xi_n$ where $\xi_n(\chi_{\bar{n}})$ is an $n(\bar{n})$-collinear field[1] and n and $\bar{n}$ are light-cone vectors satisfying $n^2 = \bar{n}^2 = 0$, $n\cdot\bar{n} = 2$. The analysis of factorization for these operators appears in Refs. [19,20] and the Wilson coefficients at next-to-leading order (NLO) in α_s have been calculated in Refs. [18,20].

The leading operator in SCET only contributes to longitudinally polarized vector meson production, but in $B \to V\gamma$ the vector meson must be transversely polarized. Transversely polarized vector mesons can be produced from subleading operators that are higher order in the SCET expansion parameter λ. The relevant effective weak Hamiltonian in SCET is

$$H_{W,\rm SCET}^{(1,4q)} = \frac{G_F}{\sqrt{2}}\sum_p\lambda_p^{(q)}\sum_{i=1}^{6}\int_0^1 dx\mathcal{B}_i^p(x,\mu)\mathcal{O}_i^{(1,4q)}(x,\mu), \quad (5)$$

where $\mathcal{O}_i^{(1,4q)}$ are

$$\begin{aligned}
\mathcal{O}_1^{(1,4q)} &= \bar{\chi}^u_{\bar{n}}W_{\bar{n}}\gamma^\perp_\mu(1-\gamma_5)Y^\dagger_{\bar{n}}b_v[\bar{\xi}^q_{\bar{n}}W_n\gamma^\mu_\perp(1-\gamma_5)W^\dagger_n\xi^u_n\\
&\quad + \bar{\xi}^q_nW_n\gamma^\mu_\perp(1-\gamma_5)W^\dagger_n\xi^u_{\bar{n}}]_x,\\
\mathcal{O}_{2,3}^{(1,4q)} &= \bar{\chi}^q_{\bar{n}}W_{\bar{n}}\gamma^\perp_\mu(1-\gamma_5)Y^\dagger_{\bar{n}}b_v[\bar{\xi}^u_{\bar{n}}W_n\gamma^\mu_\perp(1\mp\gamma_5)W^\dagger_n\xi^u_n\\
&\quad + \bar{\xi}^u_nW_n\gamma^\mu_\perp(1\mp\gamma_5)W^\dagger_n\xi^u_{\bar{n}}]_x,\\
\mathcal{O}_4^{(1,4q)} &= \sum_{q'=u,d,s}\bar{\chi}^{q'}_{\bar{n}}W_{\bar{n}}\gamma^\perp_\mu(1-\gamma_5)Y^\dagger_{\bar{n}}b_v[\bar{\xi}^q_{\bar{n}}W_n\gamma^\mu_\perp(1-\gamma_5)\\
&\quad \times W^\dagger_n\xi^{q'}_n + \bar{\xi}^q_nW_n\gamma^\mu_\perp(1-\gamma_5)W^\dagger_n\xi^{q'}_{\bar{n}}]_x,\\
\mathcal{O}_{5,6}^{(1,4q)} &= \sum_{q'=u,d,s}\bar{\chi}^q_{\bar{n}}W_{\bar{n}}\gamma^\perp_\mu(1-\gamma_5)Y^\dagger_{\bar{n}}b_v[\bar{\xi}^{q'}_{\bar{n}}W_n\gamma^\mu_\perp(1\mp\gamma_5)\\
&\quad \times W^\dagger_n\xi^{q'}_n + \bar{\xi}^{q'}_nW_n\gamma^\mu_\perp(1\mp\gamma_5)W^\dagger_n\xi^{q'}_{\bar{n}}]_x.
\end{aligned} \quad (6)$$

Here the superscript 1 denotes suppression by one power of λ compared to the leading operator, $W_{n(\bar{n})}$ is a collinear Wilson line in the $n(\bar{n})$-direction, and $Y_{\bar{n}}$ is an ultrasoft Wilson line. The subscript outside the square brackets denotes that a delta function which fixes the momentum fraction x is included in the bilinear operator:

$$[\bar{\xi}_{\bar{n}}W_n\Gamma W^\dagger_n\xi_n]_x \equiv \bar{\xi}_{\bar{n}}W_n\delta\bigg(x - \frac{\bar{\mathcal{P}}^\dagger}{2E_V}\bigg)\Gamma W^\dagger_n\xi_n, \quad (7)$$

where $\bar{\mathcal{P}} = \bar{n}\cdot\mathcal{P}$ is a derivative operator taking the largest momentum component and E_V is the energy of the produced vector meson. Here, $\xi_{\bar{n}}$ is the power-suppressed component in the spin projection of $q_n = (\not{n}\not{\bar{n}}/4)\xi_n + (\not{\bar{n}}\not{n}/4)\xi_{\bar{n}}$, where q_n is the collinear quark field. Using the equation of the motion, $\xi_{\bar{n}}$ can be expressed in terms of ξ_n

[1]Our conventions are the same as Ref. [18] except that we have exchanged $n \leftrightarrow \bar{n}$ compared to that paper.

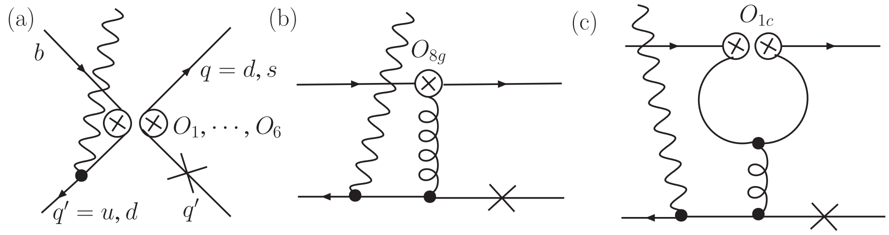

FIG. 2. Various isospin-breaking contributions in full QCD. Here crosses denote another possible photon emission from the spectator quark. If we do not consider the long-distance contribution in the diagram (c), all the contributions are power-suppressed by $\mathcal{O}(\Lambda/m_b)$.

as

$$\xi_{\bar{n}} = W_n \frac{1}{\bar{\mathcal{P}}} W_n^\dagger i\not{D}_\perp \frac{\bar{\not{n}}}{2} \xi_n. \tag{8}$$

The Wilson coefficients $\mathcal{B}_i^p$ in Eq. (5) are the same as the leading Wilson coefficients $\mathcal{C}_i^p$ of Ref. [18] due to reparametrization invariance [21]. Because the isospin-breaking contributions from $H_{W,\mathrm{SCET}}^{(1,4q)}$ are suppressed by $\Lambda_{\mathrm{QCD}}/m_b$ compared to the leading decay amplitude, we will suppress EM penguins with $C_{7,8,9,10}$ at tree level and keep only $C_{1,2,8g}$ at one-loop order in $\mathcal{B}_i^p$ for our phenomenological analysis.

The four-quark operators in the effective weak SCET Hamiltonian, Eq. (5), contribute to the radiative weak decay when the photon is emitted from the initial spectator quark, as shown in Fig. 3(a). To calculate their contribution, we need to take the time-ordered products of $\mathcal{O}_i^{(1,4q)}$ with the electromagnetic interaction term $\mathcal{L}_{\mathrm{EM}}^{(1)}$,

$$\mathcal{L}_{\mathrm{EM}}^{(1)} = e_q \bar{q}_{us} Y_{\bar{n}} \not{A}_\perp W_{\bar{n}}^\dagger \chi_{\bar{n}}^q + \mathrm{H.c.}, \tag{9}$$

where A^μ is a photon field and e_q is the electric charge of the quark. The time-ordered products are performed in $\mathrm{SCET_I}$ with possible off-shellness $p^2 \sim m_b\Lambda$ and then matched onto $\mathrm{SCET_{II}}$, which describes dynamics with fluctuations of $p^2 \sim \Lambda^2$. In the matching, $\bar{n}$-collinear fields having large off-shellness of $\mathcal{O}(m_b\Lambda)$ must be integrated out, giving a jet function

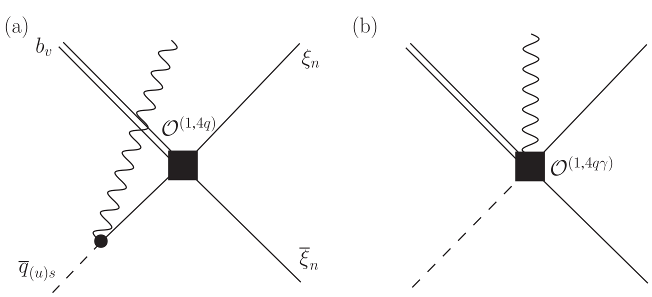

FIG. 3. SCET diagrams for the isospin-breaking corrections. Each diagram represents the electromagnetic interactions from initial and final spectator quarks, respectively.

$$\begin{aligned}\langle 0|\mathrm{T}\{W_{\bar{n}}^\dagger \chi_{\bar{n}}(z), \bar{\chi}_{\bar{n}} W_{\bar{n}}(0)\}|0\rangle = i\frac{\bar{\not{n}}}{2}\delta(z_-)\delta^2(z_\perp) \\ \times \int \frac{dk_-}{2\pi} e^{-ik_- z_+/2} J_{n\cdot p_\gamma}(k_-),\end{aligned} \tag{10}$$

where $k_+ = n \cdot k$, $k_- = \bar{n} \cdot k$, and the momentum of the photon $p_\gamma^\mu = n \cdot p_\gamma \bar{n}^\mu/2 = m_b \bar{n}^\mu/2$. At the lowest order in α_s, the jet function is simply $J_{n\cdot p_\gamma} = 1/k_-$.

The n- and $\bar{n}$-collinear degrees of freedom are decoupled at leading order in SCET, and by using a field redefinition it is possible to decouple ultrasoft degrees of freedom from both. Therefore the n-collinear piece of the matrix element describing the production of the light vector meson is decoupled from the $\bar{n}$-collinear and ultrasoft parts. Thus, we can compute B to γ from $H_{W,\mathrm{SCET}}^{(1,4q)}$, which only depends on ultrasoft and $\bar{n}$-collinear physics, independently of the light-meson production process, which depends on n-collinear physics. After a brief calculation, we find[2]

$$\begin{aligned}\hat{\mathcal{T}}_{4q}^\mu &= i\int d^4z \langle\gamma(\epsilon_\perp^*)|\mathrm{T}\{\bar{\chi}_{\bar{n}}^q W_{\bar{n}} \gamma_\mu^\perp (1-\gamma_5) \\ &\quad \times Y_{\bar{n}}^\dagger b_v(0), \mathcal{L}_{\mathrm{EM}}^{(1)}(z)\}|B\rangle, \\ &= -i\frac{e_q}{2} f_B m_B (\epsilon_\perp^{*\mu} + i\varepsilon_\perp^{\mu\nu}\epsilon_{\perp\nu}^*) \\ &\quad \times \int dl_- J_{n\cdot p_\gamma}(-l_-)\phi_B^+(l_-),\end{aligned} \tag{11}$$

where $\varepsilon_\perp^{\mu\nu} = \varepsilon^{\mu\nu\rho\sigma} n_\rho \bar{n}_\sigma/2$ setting $\varepsilon^{0123} = -1$. ϕ_B^+ is a light-cone distribution amplitude (LCDA) of the B-meson [23]. We use the convention of Ref. [24] with n and $\bar{n}$ interchanged.

The light meson production is described by the n-collinear part of the matrix elements. The matrix elements for production of transversely polarized vector mesons in SCET are

$$\begin{aligned}&\langle V_\perp(\eta_\perp^*)|[\bar{\xi}_{\bar{n}} W_n \gamma_\perp^\mu W_n^\dagger \xi_n + \bar{\xi}_n W_n \gamma_\perp^\mu W_n^\dagger \xi_{\bar{n}}]_x|0\rangle \\ &\quad = -if_V m_V \eta_\perp^{*\mu} g_\perp^{(v)}(x),\end{aligned} \tag{12}$$

[2]At higher order in α_s, this factorization theorem continues to hold but the jet function $J_{n\cdot p_\gamma}$ contains additional terms beyond what is given in Eq. (10). See Ref. [22] for details.

$$\langle V_\perp(\eta_\perp^*)|[\bar{\xi}_{\bar{n}}W_n\gamma_\perp^\mu\gamma_5W_n^\dagger\xi_n + \bar{\xi}_nW_n\gamma_\perp^\mu\gamma_5W_n^\dagger\xi_{\bar{n}}]_x|0\rangle = -\frac{f_V}{4}m_V\varepsilon_\perp^{\mu\nu}\eta_{\perp\nu}^*\frac{\partial}{\partial x}g_\perp^{(a)}(x), \tag{13}$$

where $g_\perp^{(v,a)}$ are chiral-even, twist-3 LCDAs [25] whose asymptotic forms are $g_\perp^{(v)}(x) = 3(x^2+\bar{x}^2)/2$ and $g_\perp^{(a)}(x) = 6x\bar{x}$, where $\bar{x} = 1-x$.

Combining Eqs. (11)–(13), the matrix element of $\mathcal{O}_1^{(1,4q)}$ for $B^- \to \rho^-\gamma$, for example, is

$$\begin{aligned}\langle\rho^-\gamma|\mathcal{O}_1^{(1,4q)}(x,\mu)|B^-\rangle = &-\frac{e_u}{2}f_Bf_\rho m_Bm_\rho A_L(\epsilon_\perp^*,\eta_\perp^*)\\ &\times\int dl_-J_{n\cdot p_\gamma}(-l_-,\mu,\mu_0)\\ &\times\phi_B^+(l_-,\mu_0)\bigg[g_\perp^{(v)}(x,\mu)\\ &-\frac{\partial}{4\partial x}g_\perp^{(a)}(x,\mu)\bigg],\end{aligned} \tag{14}$$

where the renormalization scales are roughly given by $\mu\sim\sqrt{m_b\Lambda_{\rm QCD}}$ and $\mu_0\sim\Lambda_{\rm QCD}$, and $A_L = \epsilon_\perp^*\cdot\eta_\perp^* - i\varepsilon_\perp^{\mu\nu}\epsilon_{\perp\mu}^*\eta_{\perp\nu}^*$ is the polarization factor for the left-handed ρ^- and γ. In the case of $B\to\rho\gamma$, it is necessary to include one-loop corrections to the jet function; see Ref. [22] for details.

When the photon radiates from a final state quark (i.e., from the crosses in Fig. 2), the intermediate quark line is hard with off-shellness of order m_b^2. Matching onto SCET we obtain localized four-quark operators with photons, shown in Fig. 3(b). In this case, the effective weak Hamiltonian that contributes to the decay amplitudes is

$$H_{W,\rm SCET}^{(1,4q\gamma)} = \frac{G_F}{\sqrt{2}}\sum_p\lambda_p^{(q)}\sum_{i=1}^2\int_0^1dx\mathcal{A}_i^p(x,\mu)\mathcal{O}_i^{(1,4q\gamma)}(x,\mu), \tag{15}$$

where the five-particle operators $\mathcal{O}_i^{(1,4q\gamma)}$ are

$$\begin{aligned}\mathcal{O}_{\{1,2\}}^{(1,4q\gamma)}(x) = &\sum_{q'=u,d,s}e_{q'}\bar{q}'Y_n\{\bar{n\!\!\!/}, n\!\!\!/\}(1+\gamma_5)Y_n^\dagger b_v\\ &\times\bigg[\bar{\xi}_{\bar{n}}^qW_n\frac{\bar{n\!\!\!/}}{2}A\!\!\!/_\perp(1+\gamma_5)\frac{1}{\bar{\mathcal{P}}}W_n^\dagger\xi_n^{q'}\bigg]_x.\end{aligned} \tag{16}$$

In Eq. (15), the Wilson coefficients $\mathcal{A}_i^p$ at NLO are

$$\begin{aligned}\mathcal{A}_1^p(x,\mu) &= C_6+\frac{C_5}{N}+\frac{\alpha_s}{4\pi}\frac{C_F}{N}\bigg\{\frac{2C_1}{3}\bigg[1+\ln\frac{m_B^2}{\mu^2}-\frac{3}{2}G(s_p,x)\bigg]-2C_{8g}\frac{m_b}{\bar{x}m_B}\bigg\},\\ \mathcal{A}_2^p(x,\mu) &= C_6+\frac{C_5}{N}+\frac{\alpha_s}{4\pi}\frac{C_F}{N}\bigg\{\frac{C_1}{3}\bigg[1+\ln\frac{m_B^2}{\mu^2}-\frac{3}{2}G(s_p,x)\bigg]-C_{8g}\frac{m_b}{m_B}\bigg\},\end{aligned} \tag{17}$$

where $G(s_p,x)$ is

$$G(s_p,x) = -4\int_0^1dzz\bar{z}\ln(s_p-z\bar{z}\bar{x}-i\epsilon), \tag{18}$$

and $s_p\equiv m_p^2/m_B^2$. Similar to $\mathcal{B}_i^p$, we neglect EM penguins at tree level and only keep $C_{1,2,8g}$ at one loop. The sum of the terms in Eq. (17) proportional to C_1 differs by a factor of 3/4 from Ref. [10].

Again, the $B\to\gamma$ piece of the matrix element factors from the light-meson production piece. The n-collinear part in Eq. (16), describing the vector meson production, gives a leading twist LCDA $\phi_\perp(x)$ whose asymptotic form is $6x\bar{x}$. It can be obtained from the following projection:

$$\begin{aligned}&\langle V_\perp(\eta_\perp^*)|\bigg[\bar{\xi}_nW_n\delta\bigg(x-\frac{\bar{\mathcal{P}}^\dagger}{2E_V}\bigg)\bigg]_a^\alpha[W_n^\dagger\xi_n]_b^\beta|0\rangle_{\text{twist}-2}\\ &\quad=-\frac{i}{2}f_V^\perp E_V\bigg(\eta\!\!\!/_\perp^*\frac{n\!\!\!/}{2}\bigg)_{ba}\frac{\delta^{\beta\alpha}}{N}\phi_\perp(x).\end{aligned} \tag{19}$$

As an example, the $B\to K^*\gamma$ matrix elements from $\mathcal{O}_i^{(1,4q\gamma)}$ are

$$\begin{aligned}\langle\bar{K}^{*0}\gamma|\mathcal{O}_1^{(1,4q\gamma)}|\bar{B}^0\rangle &= \langle\bar{K}^{*0}\gamma|\mathcal{O}_2^{(1,4q\gamma)}|\bar{B}^0\rangle\\ &= -\frac{e_d}{2}f_Bf_{K^*}^\perp m_BA_L(\epsilon_\perp^*,\eta_\perp^*)\\ &\quad\times\frac{\phi_\perp(x,\mu)}{\bar{x}}.\end{aligned} \tag{20}$$

Next we turn to the contributions from NPCP, shown in Fig. 1(a). The size of this contribution is $\mathcal{O}(v\alpha_s(2m_c))$ [15] and therefore is suppressed only logarithmically in the large m_c limit, compared to the power suppression of previously considered $\mathcal{O}(1/m_b)$ contributions. Numerically, $\Lambda_{\rm QCD}/m_b$ and $v\alpha_s(2m_c)$ are roughly the same size, so *a priori* it is sensible to include them at the same order. In fact we will see below that the NPCP gives an important contribution to isospin violation in $B\to\rho\gamma$. When $\bar{x}$ is close to $4s_c^2$, long-distance interactions govern the charm quark pair in the loop and hence it cannot be separated from the B-meson. However, the n-collinear piece of $V_\perp$ can still be decoupled, with the dominant part obtained from the leading twist projection of Eq. (19). The factorization process is similar to the treatment in Ref. [18] and we refer the reader to that paper for details. The NPCP contribution to the decay amplitude is

$$\begin{aligned}M^{c\bar{c}} &= \frac{G_F}{\sqrt{2}}\lambda_c^{(q)}\langle V_\perp\gamma|C_1O_{1c}|B\rangle_{\rm NPCP}\\ &= i\frac{G_F}{4N\sqrt{2}}\lambda_c^{(q)}f_V^\perp m_B\eta_{\perp\mu}^*\int dx\delta(\bar{x}-4s_c^2)\\ &\quad\times\phi_\perp(x)H_{c\bar{c}}(x,m_B)\langle\gamma|\mathcal{O}_{c\bar{c}}^\mu|B\rangle,\end{aligned} \tag{21}$$

where $H_{c\bar{c}} = 4C_1\pi\alpha_s/(\bar{x}m_B^2)$ at lowest order and the six-quark operator $\mathcal{O}_{c\bar{c}}^\mu$, including nonrelativistic charm quark fields $c_{\pm{\rm v}}$, is defined as

$$\mathcal{O}^{\mu}_{c\bar{c}} = i\int d^4y \bar{c}_{-v} Y_n \gamma_\perp^\nu T^a Y_{\bar{n}}^\dagger c_{+v}(y) \bar{\chi}_{\bar{n}}^{q'} W_{\bar{n}}(y) \gamma_\nu^\perp \gamma_\perp^\mu \frac{\not{n}}{2} \times \gamma^\rho (1-\gamma_5) T^a Y_n^\dagger c_{-v} \bar{c}_{+v} \gamma_\rho (1-\gamma_5) b_v(0). \quad (22)$$

In order to integrate out $\bar{n}$-collinear fields with fluctuations greater than $\Lambda_{\rm QCD}$, we consider time-ordered products of $\mathcal{O}^{\mu}_{c\bar{c}}$ with $\mathcal{L}^{(1)}_{\rm EM}$,

$$\begin{aligned}\hat{\mathcal{T}}_{c\bar{c}} &= i\eta^*_{\perp\mu}\int d^4z \langle\gamma(\epsilon^*_\perp)|\mathrm{T}\{\mathcal{O}^{\mu}_{c\bar{c}}(0), \mathcal{L}^{(1)}_{\rm EM}(z)\}|B\rangle \\ &= ie_{q'}\int d^4y \frac{dz_+ dk_-}{4\pi} e^{-ik_-(z_+-y_+)/2} J_{n\cdot p_\gamma}(k_-) \\ &\quad \times \langle 0|\mathcal{O}_{c\bar{c}q'b}(\epsilon^*_\perp, \eta^*_\perp, z_+, y)|B\rangle, \end{aligned} \quad (23)$$

where we employed Eq. (10) to obtain the second line of Eq. (23) and $\mathcal{O}_{c\bar{c}q'b}$ is

$$\begin{aligned}\mathcal{O}_{c\bar{c}q'b}(\epsilon^*_\perp, \eta^*_\perp, z_+, y) &= \overline{q'}_{us} Y_{\bar{n}}\Big(\frac{z_+}{2}, y_\perp, \frac{y_-}{2}\Big)\frac{\bar{\not{n}}}{2}\not{\epsilon}^*_\perp \gamma^\nu_\perp \not{\eta}^*_\perp \frac{\not{n}}{2} \\ &\quad \times \gamma^\rho(1-\gamma_5) T^a Y_n c_{-v}(0) \bar{c}_{-v} Y_n \gamma^\perp_\nu \\ &\quad \times T^a Y^\dagger_{\bar{n}} c_{+v}(y) \bar{c}_{+v} \gamma_\rho (1-\gamma_5) b_v(0). \end{aligned} \quad (24)$$

The matrix element of $\mathcal{O}_{c\bar{c}q'b}$ in Eq. (23) is purely nonperturbative. It can be decomposed into left- and right-handed polarized contributions,

$$\begin{aligned}&\int d^4y \langle 0|\mathcal{O}_{c\bar{c}q'b}(\epsilon^*_\perp, \eta^*_\perp, z_+, y)|B\rangle e^{ik_-y_+/2} \\ &= f_B m_B^3 \int dl_- e^{-il_-z_+/2}[A_L(\epsilon^*_\perp, \eta^*_\perp) F^L_{c\bar{c}}(k_-, l_-) \\ &\quad + A_R(\epsilon^*_\perp, \eta^*_\perp) F^R_{c\bar{c}}(k_-, l_-)], \end{aligned} \quad (25)$$

where $A_R = \epsilon^*_\perp \cdot \eta^*_\perp + i\varepsilon_\perp^{\mu\nu}\epsilon^*_{\perp\mu}\eta^*_{\perp\nu}$ is the polarization factor for decay into right-handed final states. An important point is that the NPCP can give a right-handed polarized contribution which is not $\mathcal{O}(1/m_b)$ suppressed. As pointed out in Ref. [26], any other right-handed polarized contributions to the decay amplitude should be suppressed by $1/m_b$. Therefore NPCP could give the dominant contribution to the right-handed polarized decay amplitudes. Since the right-handed polarized contribution from NPCP cannot interfere with the leading order amplitude which produces only left-handed final states, the right-handed contribution does not enter into the asymmetries until higher orders. Therefore, we neglect any possible right-handed contribution from NPCP in our calculations of the asymmetries.

Combining Eqs. (21), (23), and (25), we obtain

$$\begin{aligned}M^{c\bar{c}} &= -\frac{G_F}{\sqrt{2}}\lambda_c^{(q)} e_{q'} f_B f_V^\perp m_B^2 A_L(\epsilon^*_\perp, \eta^*_\perp)\frac{\pi\alpha_s}{N\Lambda_{c\bar{c}}} \\ &\quad \times \int_0^1 dx \frac{\phi_\perp(x)}{\bar{x}}\delta(\bar{x}-4s_c^2)\hat{H}_{c\bar{c}}(\bar{x}), \end{aligned} \quad (26)$$

where $q = d$ or s, q' is the B-meson spectator quark, and $\hat{H}_{c\bar{c}} = \bar{x} m_B^2 H_{c\bar{c}}/(4\pi\alpha_s) = C_1 + \cdots$. Here we have defined $\Lambda_{c\bar{c}}^{-1}$ to be

$$\begin{aligned}\Lambda^{-1}_{c\bar{c}} &= -\int dl_- \frac{dz_+ dk_-}{4\pi} e^{-i(k_-+l_-)z_+/2} J_{n\cdot p_\gamma}(k_-) F^L_{c\bar{c}}(k_-, l_-) \\ &= -\int dl_- J_{n\cdot p_\gamma}(-l_-) F^L_{c\bar{c}}(-l_-, l_-) \\ &\sim \int dl_- \frac{F^L_{c\bar{c}}(-l_-, l_-)}{l_-}. \end{aligned} \quad (27)$$

Following Ref. [15], we can power count the size of this correction. The NPCP contribution is suppressed relative to the leading order term by $v\alpha_s(2m_c)$, and thus of order $m_b v\alpha_s(2m_c)/\Lambda_{\rm QCD}$ relative to the other isospin-breaking terms considered. Based on this power counting, we expect that

$$\frac{m_B}{\Lambda_{c\bar{c}}}\frac{\phi_\perp(4s_c^2)}{4s_c^2} \sim \frac{v m_b}{\Lambda_{\rm QCD}}. \quad (28)$$

The factor $\phi_\perp(4s_c^2)/(4s_c^2)$ is formally $O(1)$ in the power counting of Ref. [15], but numerically $\phi_\perp(4s_c^2)/(4s_c^2) \approx 4.3$, so we keep this factor in estimating $m_B/\Lambda_{c\bar{c}}$. Taking $v m_b/\Lambda_{\rm QCD} \sim 3$ we find $m_B/\Lambda_{c\bar{c}} \sim 0.7$. The value of $m_B/\Lambda_{c\bar{c}}$ [see Eq. (47) below] we extract from fits to the isospin asymmetries $\Delta^{K^*}_{0-}$ and Δ^{ρ}_{0-}, the CP asymmetry Δ^{ρ}_{+-}, and $\mathrm{Br}[B^+ \to \rho^+\gamma]$, is consistent with this naive estimate but smaller. For these values of $m_B/\Lambda_{c\bar{c}}$, the NPCP gives significant contributions to the isospin asymmetries.

Finally, there is another interesting isospin-breaking source, a double photon contribution with the EM penguin $O_{7\gamma}$. It is only available for the decay with an unflavored vector meson, i.e., $B \to \rho^0\gamma$. The largest contributions are depicted in Fig. 4. Concentrating first on Fig. 4(a), the off-shell photon coming from $O_{7\gamma}$ produces the vector meson and then an additional photon is emitted from the B-meson spectator quark. Integrating out the hard photon, we can match onto the SCET$_{\rm I}$ operator

$$\begin{aligned}C_{7\gamma}O_{7\gamma} &\to \frac{ee_{q'}}{4\pi^2}\frac{m_b m_B}{m_V^2}\int dx \mathcal{C}_{\gamma\gamma}(x)\bar{\chi}^q_{\bar{n}} W_{\bar{n}}\frac{\not{n}}{2}\gamma^\mu_\perp(1+\gamma_5) \\ &\quad \times Y^\dagger_{\bar{n}} h_v[\bar{\xi}^{q'}_{\bar{n}} W_n \gamma^\perp_\mu W^\dagger_n \xi^{q'}_n + \bar{\xi}^{q'}_n W_n \gamma^\perp_\mu W^\dagger_n \xi^{q'}_{\bar{n}}]_x, \end{aligned} \quad (29)$$

where $C_{\gamma\gamma}$ is equal to $C_{7\gamma}$ at tree level. Next we integrate out the $\bar{n}$-collinear fields in the time-ordered product with $\mathcal{L}^{(1)}_{\rm EM}$ and match onto SCET$_{\rm II}$. Applying Eqs. (10)–(12), we find

$$\begin{aligned}\langle V_\perp\gamma|C_{7\gamma}O_{7\gamma}|B\rangle_{2\gamma(\mathrm{a})} &= \frac{ee_q e_{q'}}{8\pi^2} f_B f_V \frac{m_b m_B^2}{m_V} C_{7\gamma} A_L(\epsilon^*_\perp, \eta^*_\perp) \\ &\quad \times \int dl_- J_{n\cdot p_\gamma}(-l_-)\phi^+_B(l_-), \end{aligned} \quad (30)$$

and, in the case of $B \to \rho^0\gamma$,

$$\langle \rho^0_\perp \gamma | C_{7\gamma} O_{7\gamma} | \bar{B}^0 \rangle_{2\gamma(\mathrm{a})} = \frac{Q_d(Q_u - Q_d)}{\sqrt{2}} \frac{e\alpha}{2\pi} f_B f_\rho \times \frac{m_b m_B^2}{m_\rho} C_{7\gamma} A_L(\epsilon^*_\perp, \eta^*_\perp) \times \int dl_- J_{n\cdot p_\gamma}(-l_-)\phi_B^+(l_-), \quad (31)$$

where $Q_q = e_q/e$ with $Q_u = 2/3$ and $Q_d = -1/3$.

At the lowest order in α_s, the contribution of Fig. 4(b) is the same as Fig. 4(a) with n and $\bar{n}$ exchanged. So the result can be written as

$$\langle \rho^0_\perp \gamma | C_{7\gamma} O_{7\gamma} | \bar{B}^0 \rangle_{2\gamma(\mathrm{b})} = \frac{Q_d(Q_u - Q_d)}{\sqrt{2}} \frac{e\alpha}{2\pi} f_B f_\rho \times \frac{m_b m_B^2}{m_\rho} C_{7\gamma} A_L(\epsilon^*_\perp, \eta^*_\perp) \times \int dl_+ J_{n\cdot p_\gamma}(-l_+)\phi_B^+(l_+). \quad (32)$$

The double-photon contribution is suppressed by α, but enhanced by a factor of m_b^2/m_V^2 due to the virtual photon. Compared to the other isospin-breaking contributions, such as those in Eqs. (14) and (20), this contribution is rather small.

The isospin asymmetries in Eq. (1) are given by

$$\Delta^V_{0-} = \frac{\mathrm{Re}(b^V_d - b^V_u) + R\,\mathrm{Re}(\bar{b}^V_d - \bar{b}^V_u)}{1+R}, \quad (33)$$

where

$$b^V_d = \frac{A^V_0}{c_V L_V}, \qquad b^V_u = \frac{A^V_-}{L_V}, \quad (34)$$

$A^V_{0,-}$ are the leading isospin-breaking corrections to the decay amplitude, L_V are the leading isospin symmetric decay amplitudes, and $c_V = 1$ for K^*, $c_V = -1/\sqrt{2}$ for ρ. In Eq. (33), $\bar{b}^V_{u,d}$ are the corresponding ratios for the charge conjugate modes, and $R = |\overline{L_V}|^2/|L_V|^2$. L_V can be written as

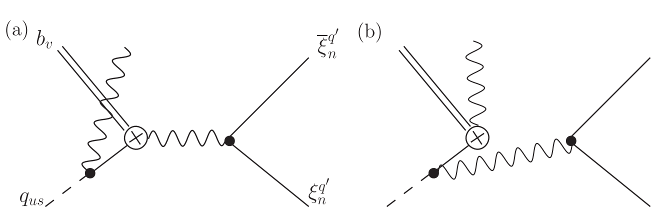

FIG. 4. Leading double photon contribution to the isospin asymmetry, where ⊗ represents $O_{7\gamma}$. They are suppressed by $\alpha m_b^2/m_V^2$ compared to other usual isospin-breaking contributions.

$$L_{K^*} = \frac{G_F}{\sqrt{2}} \frac{e}{4\pi^2} m_b^2 m_B A_L(\epsilon^*_\perp, \eta^*_\perp)\lambda_c^{(s)} a^c_{7,K^*}\zeta^{K^*}_\perp,$$
$$L_\rho = \frac{G_F}{\sqrt{2}} \frac{e}{4\pi^2} m_b^2 m_B A_L(\epsilon^*_\perp, \eta^*_\perp) \sum_{p=u,c} \lambda_p^{(d)} a^p_{7,\rho}\zeta^\rho_\perp, \quad (35)$$

where the transition form factor for $B \to V$, $\zeta^V_\perp$, is defined as

$$\langle V(\eta^*_\perp)|\bar{\xi}_n W_n \gamma^\mu_\perp (1-\gamma_5) Y_n^\dagger b_v |B\rangle = m_B(i\varepsilon_\perp^{\mu\nu}\eta^*_{\perp\nu} - \eta^{*\mu}_\perp)\zeta^V_\perp, \quad (36)$$

and we will use $\zeta^{K^*}_\perp = 0.36 \pm 0.07$ and $\zeta^\rho_\perp = 0.27 \pm 0.05$ [11] for the numerical analysis. The Wilson coefficients $a^p_{7,V}$ in Eq. (35) are

$$a^p_{7,V} = C_{7\gamma}A_7^{(0)} + \frac{\alpha_s C_F}{4\pi}[C_1 G_1(s_p) + C_{8g}G_8(s_p)] + \pi\alpha_s \frac{C_F}{N} \frac{f_B f_V^\perp m_B}{m_b^2} \int dl_+ \frac{\phi_B^+(l_+)}{l_+} \int_0^1 dx \frac{\phi_\perp(x)}{\bar{x}} \times \left[C_{7\gamma} + \frac{C_1}{6}H(x, s_p) + \frac{C_{8g}}{3}\right]\frac{1}{\zeta^V_\perp}, \quad (37)$$

where the hard functions $A_7^{(0)}$, $G_{1,8}$, and H are available in Refs. [3,5,27], and we followed the conventions of Ref. [5].

Finally we obtain

$$b^{K^*}_q = Q_q \frac{2\pi^2 f_B}{m_b a^c_{7,K^*}\zeta^{K^*}_\perp}\left[2\frac{f^{K^*}_\perp}{m_b}K^{K^*}_1 + \frac{f_{K^*}m_{K^*}}{\lambda_B m_b}K^{K^*}_{2q}\right], \quad (38)$$

$$b^\rho_q = Q_q \frac{2\pi^2 f_B}{m_b \sum_{p=u,c}\lambda_p^{(d)} a^p_{7,\rho}\zeta^\rho_\perp}\left[2\frac{f^\rho_\perp}{m_b}K^\rho_1 + \frac{f_\rho m_\rho}{\lambda_B m_b}K^\rho_{2q}\right]. \quad (39)$$

Here $\lambda_B^{-1} = \int dl \phi_B^+(l)/l$ and we use the following model for $\phi_B^+(l)$ [28]:

$$\phi_B^+(l, \mu) = \frac{4\mu l}{\pi\lambda_B(l^2+\mu^2)}\left[\frac{\mu^2}{l^2+\mu^2} - \frac{2(\sigma_B - 1)}{\pi^2}\ln\frac{l}{\mu}\right], \quad (40)$$

where the parameters λ_B and σ_B are $\lambda_B = 460 \pm 110$ MeV, $\sigma_B = 1.4 \pm 0.4$ at $\mu = 1$ GeV. $K_{1,2}$ can be written as

$$K^{K^*}_1 = \int_0^1 dx \frac{\phi_\perp(x)}{\bar{x}}\left\{-\frac{1}{2}[\mathcal{A}^c_1(x) + \mathcal{A}^c_2(x)] - C_1\frac{\pi\alpha_s}{N}\frac{m_B}{\Lambda_{c\bar{c}}}\delta(\bar{x} - 4s_c^2)\right\}, \quad (41)$$

$$K^{K^*}_{2q} = \int_0^1 dx\left[g^{(v)}_\perp - \frac{\partial}{4\partial x}g^{(a)}_\perp\right](x)\left\{\frac{\lambda_u^{(s)}}{\lambda_c^{(s)}}\left(C_1 + \frac{C_2}{N}\right)\delta_{qu} + \mathcal{B}^c_4(x, m_b)\right\}, \quad (42)$$

$$K_1^\rho = \sum_{p=u,c} \lambda_p^{(d)} \int_0^1 dx \frac{\phi_\perp(x)}{\bar{x}} \bigg\{ -\frac{1}{2}[\mathcal{A}_1^p(x) + \mathcal{A}_2^p(x)] - \delta_{pc} C_1 \frac{\pi\alpha_s}{N} \frac{m_B}{\Lambda_{c\bar{c}}} \delta(\bar{x} - 4s_c^2) \bigg\}, \tag{43}$$

$$K_{2q}^\rho = \sum_{p=u,c} \lambda_p^{(d)} \bigg\{ \int_0^1 dx \bigg[g_\perp^{(v)} - \frac{\partial}{4\partial x} g_\perp^{(a)} \bigg](x) \times \bigg[-\lambda_B \int dl_- \phi_B^+(l_-) J_{p_\gamma}(-l_-)(\delta_{qu}\mathcal{B}_1^p(x) - \delta_{qd}\mathcal{B}_2^p(x)) + \mathcal{B}_4^p(x) \bigg] + 2\delta_{qd} C_{7\gamma} \frac{\alpha}{\pi} \frac{m_b m_B}{m_\rho^2} \bigg\}. \tag{44}$$

Here we include only the tree-level contributions to Cabibbo-suppressed terms with $\lambda_u^{(s)}$ in Eq. (42) because it is numerically comparable to the other term with $\mathcal{B}_4^c$.

In the convolutions of $\mathcal{A}_1$ and $\phi_\perp(x)/\bar{x}$ in Eqs. (41) and (43) and $\mathcal{B}_4$ and $g_\perp^{(v)}$ in Eqs. (42) and (44), there are endpoint divergences, which can be eliminated with the zero-bin subtractions [9]

$$\int_0^1 dx \frac{\phi_\perp(x)}{\bar{x}^2} \to \int_0^1 dx \frac{\phi_\perp(x) + \bar{x}\phi'_\perp(1)}{\bar{x}^2}, \qquad \int_0^1 dx \frac{g_\perp^{(v)}(x)}{\bar{x}} \to \int_0^1 dx \frac{g_\perp^{(v)}(x) - g_\perp^{(v)}(1)}{\bar{x}}. \tag{45}$$

The zero-bin subtraction removes infrared divergences from the x integrals. We have dropped all finite terms including logarithms associated with rapidity scale dependence. We estimate the uncertainty associated with this procedure to be 50%.

For numerical estimates of $b_q^{K^*,\rho}$ in Eqs. (38) and (39), we use the following set of parameters: $\{m_b, m_c, m_B, m_{K^*}, m_\rho, f_B, f_{K^*}, f_\rho, f_\perp^{K^*}, f_\perp^\rho\} = \{4.8, 1.3, 5.28, 0.894, 0.775, 0.2 \pm 0.03, 0.218, 0.209, 0.175 \pm 0.025, 0.150 \pm 0.025\}$ GeV. The CKM parameters are $\bar{\rho} = 0.221 \pm 0.064$ and $\bar{\eta} = 0.340 \pm 0.045$. All Wilson coefficients and hard functions are evaluated at the scale $\mu = m_b$, we do not include any renormalization group evolution, and we use the asymptotic forms for the vector meson wave function $\phi_\perp$ and $g_\perp^{(v),(a)}$. Our estimates for the isospin asymmetries in the absence of NPCP contributions are

$$\Delta_{0-}^{K^*} = 0.04 \pm 0.02, \qquad \Delta_{0-}^{\rho} = 0.02 \pm 0.02, \tag{46}$$

where the dominant errors come from λ_B, $\zeta_\perp^V$, and CKM factors. These estimates for Δ_{0-}^V are comparable to previous theoretical results [10,11,29]. Comparing to Eq. (2), we see that $\Delta_{0-}^{K^*}$ is consistent, but Δ_{0-}^{ρ} disagrees by about 1.7 σ.

Next we include the NPCP contribution in our calculation. In addition to the isospin asymmetries, the NPCP can contribute to the CP-violating asymmetry, Δ_{+-}^{ρ} [17], and to the branching ratio, $\mathrm{Br}[B^+ \to \rho^+\gamma]$ [2]. In order to obtain values of $m_B/\Lambda_{c\bar{c}}$ that are not inconsistent with measurements of these quantities, we perform a least squares fit to all four observables. In our calculations of Δ_{+-}^{ρ} and $\mathrm{Br}[B^+ \to \rho^+\gamma]$, we include only the leading order and NPCP contributions, without any $O(1/m_b)$ corrections. An analysis that includes the $O(1/m_b)$ corrections to all four observables is clearly required but is beyond the scope of this paper.

The results of the fits along with experimental results are shown in Table I. The first column lists the observables considered and the second gives their measured values including errors. In the third column, we show the theoretical prediction in the absence of the NPCP contribution for the values of the parameters given earlier. The last column gives the results of the fit with the NPCP included. We extract

$$\mathrm{Re}\bigg[\frac{m_B}{\Lambda_{c\bar{c}}}\bigg] = -0.102 \pm 0.063, \qquad \mathrm{Im}\bigg[\frac{m_B}{\Lambda_{c\bar{c}}}\bigg] = 0.022 \pm 0.255. \tag{47}$$

The χ^2 of the predictions in column three of Table I is 15.2. Including the NPCP, the χ^2 is 12.1, so the overall agreement between experiment and theory is slightly improved. Note that after including the NPCP the theoretical prediction for Δ_{0-}^{ρ} increases so that the 1σ error band of the experimental result and the theoretical result now overlap. However, the prediction for $\Delta_{0-}^{K^*}$ is now significantly increased. The trend suggested by the central values of the experimental data, namely, a large value of Δ_{0-}^{ρ} and small value of $\Delta_{0-}^{K^*}$, does not seem to be naturally accommodated by including NPCP contributions. However, once theoretical and experimental uncertainties are taken into account, the theoretical predictions are consistent with both isospin asymmetries. For the range of values of $m_B/\Lambda_{c\bar{c}}$ obtained in our fit, the NPCP does not have significant impact on the theoretical predictions for Δ_{+-}^{ρ} and $\mathrm{Br}[B^+ \to \rho^+\gamma]$. Finally, inclusion of NPCP contributions substantially increases the uncertainty in all theoretical predictions because the parameter $m_B/\Lambda_{c\bar{c}}$ is not well constrained.

These results indicate that the NPCP can increase the isospin-violating asymmetry Δ_{0-}^{ρ} to bring theoretical predictions closer to current data, while maintaining consistency with the other observed asymmetries. It is not possible to obtain predictions for the two isospin asymmetries that are in agreement with the central values of both $\Delta_{0-}^{K^*}$ and Δ_{0-}^{ρ}. However, the uncertainties in the current measurements of all asymmetries are large and better measurements are needed to determine whether the NPCP is an important contribution to $B \to \rho\gamma$ isospin and CP-violating asymmetries.

To summarize, we have used SCET to calculate the isospin asymmetries in $B \to V\gamma$ decays, including all $\mathcal{O}(1/m_b)$ contributions as well as the $\mathcal{O}(v\alpha_s)$ NPCP con-

TABLE I. Theoretical predictions with and without NPCP compared to experimental data.

	Experimental results	Without NPCP	With NPCP
$\Delta_{0-}^{K^*}$	0.03 ± 0.04	0.04 ± 0.02	0.10 ± 0.05
Δ_{0-}^{ρ}	0.26 ± 0.14	0.02 ± 0.02	0.10 ± 0.06
Δ_{+-}^{ρ}	0.11 ± 0.33	0.08 ± 0.02	0.07 ± 0.13
$\mathrm{Br}[B^+ \to \rho^+ \gamma] \times 10^6$	0.96 ± 0.24	1.80 ± 0.69	1.63 ± 0.67

tribution. As in nonleptonic B-decays [30], our analysis allows for large NPCP contributions, which could account for the large isospin asymmetries measured (with large errors) in $B \to \rho\gamma$. If the isospin asymmetries are large and the NPCP is the source of these asymmetries, then we also expect the NPCP to contribute to CP asymmetries in $B \to \rho\gamma$ and give larger than expected contributions to the right-handed polarized decay rates in $B \to V\gamma$. We speculate that the NPCP could also be responsible for the recently measured enhanced transversely polarized decay amplitude for $B \to VV$ [31]. More precise experiments will be needed to determine the exact size of NPCP and confirm these predictions.

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