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Annihilation contributions in $B \to K_1\gamma$ decay in next-to-leading order in LEET and CP-asymmetry

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Abstract. The effect of weak annihilation and u-quark penguin contribution on the branching ratio $B \to K_1\gamma$ at next-to-leading order of α_s are calculated using the LEET approach. It is shown that the value of the LEET form factor remains the same in the range of the unitarity triangle phase α favored by the standard model. The CP-asymmetry for the above mentioned decay has been calculated, and its suppression due to the hard-spectator correction has also been incorporated. In addition, the sensitivity of the CP-asymmetry on the underlying parameters has been discussed.

Exclusive decays involving the $b \to s\gamma$ transition are best exemplified by the decay $B \to K^*\gamma$, which abundantly provide an issue for both theorists and experimentalists. Higher resonances of kaons such as $K_2^*(1430)$ are also measured by CLEO [1] and the B factories [2, 3]. Recently, Belle [4] has announced the first measurement of $B \to K_1^+(1270)\gamma$,

$$\mathcal{B}(B^+ \to K_1^+\gamma) = (4.28 \pm 0.94 \pm 0.43) \times 10^{-5}\,. \quad (1)$$

There are several reason to focus on higher kaon resonances. The first and most promising is that they share many features with $B \to K^*\gamma$, like that at the quark level both of them are governed by $b \to s\gamma$. Therefore all the achievements of $b \to s\gamma$ can be used in these decays, e.g. the same operators in the operator product expansion and the same Wilson coefficients that are available. The light-cone distribution amplitudes (DA) are the same except for the overall factor of γ_5, and this gives a few differences in many calculations [5–7]. Secondly, it was suggested that $B \to K_{\rm res}\,(\to K\pi\pi)\,\gamma$ can provide a direct measurement of the photon polarization [8, 9], and it was shown that a large polarization asymmetry of $\approx 33\%$ is produced due to decay of the B-meson through the kaon resonances. In the presence of anomalous right-handed couplings, the polarization can be severely reduced in the parameter space allowed by current experimental bounds of $B \to X_s\gamma$. It was also argued that the B factories can now make a lot of $B\bar{B}$ pairs, enough to check the anomalous couplings through the measurement of the photon polarization.

Theorists also face challenges from the discrepancy between their predictions and experiments. It was pointed out that the form factor obtained using the LEET approach for $B \to K^*\gamma$ is smaller compared to the values obtained by QCD sum rules or light-cone sum rules (LCSR) [10]. At this stage, the source of this mismatch is not well understood.

For the $B \to K_1\gamma$ decay the situation is more complicated. Based on the QCDF framework combined with the LCSR results, it is predicted that $\mathcal{B}(B^0 \to K_1^0(1270)\gamma) = (0.828 \pm 0.335) \times 10^{-5}$ at the NLO of α_s, which is very small as compared to the experimental value [cf. (1)] [5–7]. The value of the relevant form factor has been extracted from the experimental data and is found to be $F_+^{K_1(1270)}(0) = 0.32 \pm 0.03$, which is large as compared to $F_+^{K_1(1270)}(0)|_{\rm LCSR} = 0.14 \pm 0.03$ as obtained by the LCSR. This is contrary to the case of $B \to K^*\gamma$, where the form factor obtained from LCSR is larger than the LEET value, and the source of the discrepancy is not yet known. For the $B \to K_1\gamma$ case the possible candidates to explain this discrepancy, like higher twist effects in DA, non-zero mass effects of the axial kaon, the framework of QCDF, possible mixing in $K_1\,(1270)$ and $K_1\,(1400)$ and annihilation topologies, have also been discussed in detail in the literature [11]. The calculation done in [11] is for the leading twist, and it was pointed out that a higher twist may have some effect on the form factors, because all others are no suitable candidates. Recently it has been shown that the value of the form factor is not sensitive to the higher twists [12].

In this paper the effect of weak annihilation and also the u-quark contribution A^u from the penguin to the branching ratio for $B \to K_1\gamma$ at NLO of α_s are calculated using the LEET approach [16, 17]. We have followed the same procedure as in the work of Ali et al. [10] for $B \to K^*\gamma$, because $B \to K_1\gamma$ shares many of its characteristics with it. As it is pointed out in the literature on $B \to K^*\gamma$, the effect of the annihilation contribution to the charmed quark part of the amplitude is numerically small, because only

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the penguin operator with tiny Wilson coefficients can contribute. On the other hand the annihilation contribution to the up-quark part of the amplitude contributes significantly, because of the large Wilson coefficients but again the CKM suppression $\left|\lambda_u^{(s)}/\lambda_c^{(s)}\right| \approx 0.02$ puts this large correction for $B \to K_1\gamma$ into perspective [18]. Finally, by incorporating these annihilation and u-quark contributions we compute the CP-asymmetry $\mathcal{A}_{\text{CP}}\left(K_1^{\pm}\gamma\right)$ involving the decay $B \to K_1\gamma$. The CP-asymmetry arises due to the interference of the various penguin amplitudes which have clashing weak phases, with the required strong interaction phase provided by the $\mathcal{O}(\alpha_s)$ corrections entering the penguin amplitudes via the Bander–Silverman–Soni (BSS) mechanism [19, 20]. We find that the hard-spectator corrections reduce the CP-asymmetry calculated from the vertex contributions alone. The resulting CP-asymmetry depends rather sensitively on the ratio of the quark masses m_c/m_b. This parametric dependence, combined with the scale dependence of $\mathcal{A}_{\text{CP}}\left(K_1^{\pm}\gamma\right)$, makes the prediction of direct CP-asymmetry rather unreliable and the present work will be devoted to the study of this issue.

The effective Hamiltonian for $b \to s\gamma$ can be written as

$$\mathcal{H}_{\text{eff}}(b \to s\gamma) = -\frac{G_F}{\sqrt{2}} V_{tb}V_{ts}^* \sum_{i=1}^{8} C_i(\mu)O_i(\mu)\,, \qquad (2)$$

where

$$\begin{aligned}
O_1^{(p)} &= (\bar{s}_i p_j)_{V-A}(\bar{p}_j b_i)_{V-A}\,,\\
O_2^{(p)} &= (\bar{s}_i p_i)_{V-A}(\bar{p}_j b_j)_{V-A}\,,\\
O_3 &= (\bar{s}_i b_i)_{V-A}\sum_{q=u,c,t}(\bar{q}_j q_j)_{V-A}\,,\\
O_4 &= (\bar{s}_i b_j)_{V-A}\sum_{q=u,c,t}(\bar{q}_j q_i)_{V-A}\,,\\
O_5 &= (\bar{s}_i b_i)_{V-A}\sum_{q=u,c,t}(\bar{q}_j q_j)_{V+A}\,,\\
O_6 &= (\bar{s}_i b_j)_{V-A}\sum_{q=u,c,t}(\bar{q}_j q_i)_{V+A}\,,\\
O_7 &= \frac{em_b}{8\pi^2}\bar{s}_i\sigma^{\mu\nu}(1+\gamma_5)b_i F_{\mu\nu}\,,\\
O_8 &= \frac{g_s m_b}{8\pi^2}\bar{s}_i\sigma^{\mu\nu}(1+\gamma_5)T_{ij}^a b_j G_{\mu\nu}^a\,. \qquad (3)
\end{aligned}$$

Here i, j are color indices and p stands for the u- or c-quark, and G_F is the Fermi coupling constant. We neglect the CKM element $V_{ub}V_{us}^*$ as well as the s-quark mass. The leading contribution to $B \to K_1\gamma$ comes from the electromagnetic operator O_7 as shown in Fig. 1.

As in the case of real photon emission ($q^2 = 0$), the only form factor appearing in the calculation is $\xi_\perp^{(K_1)}$. Therefore one can write

$$\begin{aligned}
\langle O_7\rangle_A &\equiv \langle K_1(p',\epsilon)\gamma(q,e)|O_7|B(p)\rangle\\
&= \frac{em_b}{4\pi^2}\xi_\perp^{(K_1)}\left[\epsilon^* q(p+p')e^* - \epsilon^* e^*(p^2 - p'^2)\right.\\
&\quad\left. + \mathrm{i}\epsilon_{\mu\nu\alpha\beta}e^{*\mu}\epsilon^{*\nu}q^\alpha(p+p')^\beta\right]\,, \qquad (4)
\end{aligned}$$

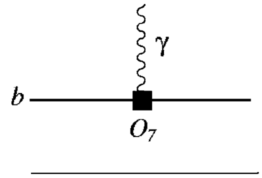

Fig. 1. Leading order contribution by the operator O_7

with $\epsilon^{*\nu}$ and e^μ being the polarization vectors for axial kaon and photon, respectively. The decay rate is straightforwardly obtained to be [5–7]

$$\begin{aligned}
\Gamma(B \to K_1\gamma) &= \frac{G_F^2 \alpha m_b^2 m_B^3}{32\pi^4}|V_{tb}V_{ts}^*|^2\\
&\quad\times\left(1 - \frac{m^2}{m_B^2}\right)^3\left|\xi_\perp^{(K_1)}\right|^2\left|C_7^{\text{eff}(0)}\right|^2\,, \qquad (5)
\end{aligned}$$

where α is the fine-structure constant, $\alpha = \alpha(0) = 1/137$ and $C_7^{\text{eff}(0)}$ is the effective Wilson coefficient at leading order.

At next-to-leading order of α_s, one has to consider the contributions from the operators O_2 and O_8 along with that of O_7 in $B \to K_1\gamma$ decay. For the operator O_7 all the subleading contributions shown in Fig. 2 are absorbed in the form factor, whereas the Wilson coefficient contains next-to-leading order parts

$$C_7^{\text{eff}}(\mu) = C_7^{\text{eff}(0)}(\mu) + \frac{\alpha_s(\mu)}{4\pi}C_7^{\text{eff}(1)}(\mu)\,.$$

On the other hand, for operators O_2 and O_8 the leading order $C_2^{(0)}$ and $C_8^{(0)}$ are sufficient because these operators contribute at NLO. Each operator has its vertex contribution and hard-spectator contribution terms which are calculated explicitly in [12] and are depicted in Figs. 2–6

The branching ratio for $B \to K_1\gamma$ is given by

$$\begin{aligned}
\mathcal{B}_{\text{th}}(B \to K_1\gamma) &= \tau_B\,\Gamma_{\text{th}}(B \to K_1\gamma)\\
&= \tau_B\,\frac{G_F^2\alpha|V_{tb}V_{ts}^*|^2}{32\pi^4}\,m_{b,\text{pole}}^2 M^3\left[\xi_\perp^{(K_1)}\right]^2\\
&\quad\times\left(1 - \frac{m_{K_1}^2}{M^2}\right)^3\left|C_7^{(0)\text{eff}} + A^{(1)}(\mu)\right|^2\,, \qquad (6)
\end{aligned}$$

where $m_{b,\text{pole}}$ is the pole b-quark mass, M and m_{K_1} are the B- and K_1-meson masses, and τ_B is the lifetime of

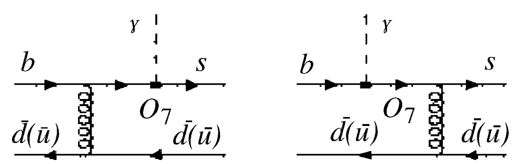

Fig. 2. Feynman diagram contributing to the spectator corrections involving the O_7 operator in the decay $B \to K_1\gamma$. The *curly* (*dashed*) *line* here and in the subsequent figures represents a gluon (photon)

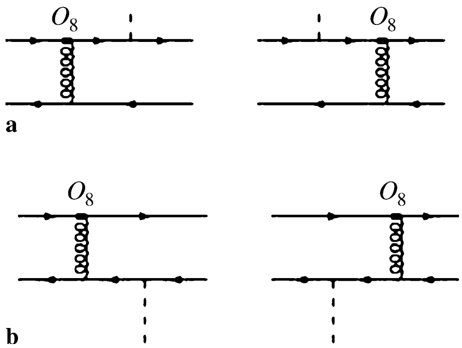

Fig. 3. Feynman diagram contributing to the spectator corrections involving the O_8 operator in the decay $B \to K_1\gamma$. *Row a*: a photon is emitted from the flavor-changing line; *row b*: photon radiation off the spectator quark line

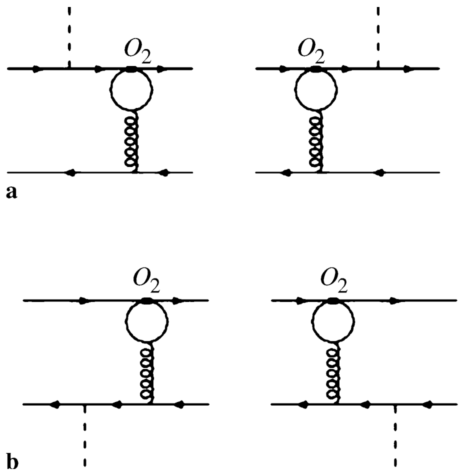

Fig. 4. Feynman diagram contributing to the spectator corrections involving the O_2 operator in the decay $B \to K_1\gamma$. *Row a*: a photon is emitted from the flavor-changing line; *row b*: photon radiation off the spectator quark line

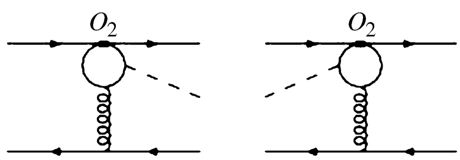

Fig. 5. Feynman diagram contributing to the spectator corrections involving the O_2 operator for the case when both the photon and virtual gluon are emitted from the internal (loop) quark line

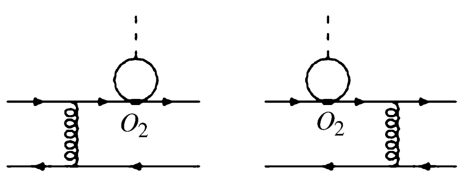

Fig. 6. Feynman diagram contributing to the spectator corrections involving the O_2 operator for the case when only the photon is emitted from the internal (loop) quark line in the $bs\gamma$ vertex

the B^0- or B^+-meson. The values of these constants are taken from [10] for the numerical analysis. For this study, we consider $\xi_\perp^{(K_1)}$ as a free parameter, and we will extract its value from the current experimental data on $B \to K_1\gamma$ decays.

The function $A^{(1)}$ in (6) can be decomposed into the following three components:

$$A^{(1)}(\mu) = A^{(1)}_{C_7}(\mu) + A^{(1)}_{\text{ver}}(\mu) + A^{(1)K_1}_{\text{sp}}(\mu_{\text{sp}})\,. \qquad (7)$$

Here, $A^{(1)}_{C_7}$ and $A^{(1)}_{\text{ver}}$ are the $O(\alpha_\text{s})$ (i.e. NLO) corrections due to the Wilson coefficient C_7^{eff} and in the $b \to s\gamma$ vertex, respectively, and $A^{(1)K_1}_{\text{sp}}$ is the $\mathcal{O}(\alpha_\text{s})$ hard-spectator correction to the $B \to K_1\gamma$ amplitude. Their explicit expressions are as follows:

$$A^{(1)}_{C_7}(\mu) = \frac{\alpha_\text{s}(\mu)}{4\pi}\, C_7^{(1)\text{eff}}(\mu)\,, \qquad (8)$$

$$\begin{aligned} A^{(1)}_{\text{ver}}(\mu) = \frac{\alpha_\text{s}(\mu)}{4\pi}\Bigg\{ &\frac{32}{81}\Big[13C_2^{(0)}(\mu) + 27C_7^{(0)\text{eff}}(\mu) \\ &-9C_8^{(0)\text{eff}}(\mu)\Big] \ln\frac{m_b}{\mu}\, \frac{20}{3}\, C_7^{(0)\text{eff}}(\mu) \\ &- + \frac{4}{27}\left(33 - 2\pi^2 + 6\pi \text{i}\right) C_8^{(0)\text{eff}}(\mu) \\ &+ r_2(z)\, C_2^{(0)}(\mu) \Bigg\}\,, \end{aligned} \qquad (9)$$

$$\begin{aligned} A^{(1)K_1}_{\text{sp}}(\mu_{\text{sp}}) = \frac{\alpha_\text{s}(\mu_{\text{sp}})}{4\pi}\, \frac{2\Delta F_\perp^{(K_1)}(\mu_{\text{sp}})}{9\xi_\perp^{(K_1)}} \Bigg\{ &3C_7^{(0)\text{eff}}(\mu_{\text{sp}}) \\ + C_8^{(0)\text{eff}}(\mu_{\text{sp}}) &\left[1 - \frac{6a_{\perp 1}^{(K_1)}(\mu_{\text{sp}})}{\langle \bar{u}^{-1}\rangle_\perp^{(K_1)}(\mu_{\text{sp}})}\right] \\ + C_2^{(0)}(\mu_{\text{sp}}) &\left[1 - \frac{h^{(K_1)}(z, \mu_{\text{sp}})}{\langle \bar{u}^{-1}\rangle_\perp^{(K_1)}(\mu_{\text{sp}})}\right] \Bigg\}\,. \end{aligned} \qquad (10)$$

The terms proportional to $\Delta F_\perp^{(\rho)}(\mu_{\text{sp}})$ above are the $O(\alpha_\text{s})$ hard-spectator corrections which should be evaluated at the typical scale $\mu_{\text{sp}} = \sqrt{\mu \Lambda_\text{H}}$ of the gluon virtuality. The complex function $r_2(z)$ of the parameter $z = m_c^2/m_b^2$, and the Wilson coefficients in the above equations can be found in [13–15]; the function $h^{(\rho)}(z, \mu)$ and the dimensionless quantity $\Delta F_\perp^{(\rho)}(\mu)$ are defined through (25) and (27), respectively, of [12]. Now $C_7^{(1)\text{eff}}(\mu)$ and $A^{(1)}_{\text{ver}}(\mu)$ are process independent and encode the QCD effects only, whereas $A^{(1)}_{\text{sp}}(\mu_{\text{sp}})$ contains the key information about the out-going mesons. The factor $\frac{6a_{\perp 1}^{(K_1)}(\mu_{\text{sp}})}{\langle \bar{u}^{-1}\rangle_\perp^{(K_1)}(\mu_{\text{sp}})}$ appearing in (10) arises due to the Gegenbauer moments.

By calculating the numerical value from the above expressions and varying the parameters in the standard range, the value of the form factor is extracted from the experimental measurements (1) and is found to be [11]

$$\xi_\perp^{(K_1)}(0) = 0.32 \pm 0.03\,,$$

which is for the leading twist and remains unchanged if one includes the higher twist effects [12].

It is already pointed out in the literature that it is unlikely that the annihilation topology would give considerable contributions [11], but these are important if one wants to study the CP-asymmetry and this is one of the purposes of this article. Before calculating the CP-asymmetry we will check the effects of the annihilation contribution on the branching ratio of $B \to K_1\gamma$ decays.

Since weak annihilation is a power correction, we will content ourselves with the lowest order result ($O(\alpha_s^0)$) for our estimate with a check of its effect on the branching ratio. The reason for including this class of power corrections is that they come with numerical enhancement from the large Wilson coefficients $C_{1,2}$ ($C_1 \approx 3C_7$) but are CKM suppressed, and thus these contributions are expected to be very small for the decay under consideration. The amplitude for charged B-meson decay in terms of weak annihilation A, charmed penguin P_c, gluonic penguin M and short-distance amplitude P_t can be written as [following the notation of [25]]

$$A\left(B^- \to K_1^-\gamma\right) = \lambda_u^{(s)} a + \lambda_t^{(s)} p\,, \tag{11}$$

$$A\left(B^0 \to K_1^0\gamma\right) = \lambda_t^{(s)}\left(P_t + \left(M^{(1)} - P_c^{(1)}\right) + \frac{2}{3}\left(M^{(2)} - P_c^{(2)}\right)\right), \tag{12}$$

where $\lambda_q^{(s)} = V_{qb}V_{qs}^*$, $a = A - P_c$ and $p = P_t + M - P_c$. As it is known [25]

$$P_c \simeq 0.2A, A \simeq 0.3P_t\,,$$

i.e. we can safely neglect charmed penguin P_c and gluonic penguin M amplitudes relative to the short-distance amplitude P_t and the weak annihilation amplitude A. Thus (11) becomes

$$A\left(B^- \to K_1^-\gamma\right) = \lambda_t^{(s)} p\left(1 + \frac{\lambda_u^{(s)}}{\lambda_t^{(s)}}\frac{a}{p}\right) = \lambda_t^{(s)} p\left(1 + \epsilon_A e^{i\phi_A}\frac{\lambda_u^{(s)}}{\lambda_t^{(s)}}\right)$$

and

$$A\left(B^0 \to K_1^0\gamma\right) = \lambda_t^{(s)} p\,,$$

where $\epsilon_A e^{i\phi_A} \equiv a/p$, ϕ_A is the strong interaction phase which disappears in $\mathcal{O}(\alpha_s)$ in the chiral limit. Hence we will set it equal to zero in the subsequent calculation. Following the same lines as for the charged B-meson, the ratio of the branching ratios for charged to neutral B-meson decays can be written as

$$\frac{\mathcal{B}\left(B^- \to K_1^-\gamma\right)}{\mathcal{B}\left(B^0 \to K_1^0\gamma\right)} \simeq \left|1 + \epsilon_A e^{i\phi_A}\frac{V_{ub}V_{us}^*}{V_{tb}V_{ts}^*}\right|^2 . \tag{13}$$

The estimates in the frame work of the light-cone QCD sum rules yield typically [26, 27] $\epsilon_A = -0.35$ and $\epsilon_A = 0.046$ for the decays $B^- \to K_1^-\gamma$ and $B^0 \to K_1^0\gamma$, respectively. Let us define

$$\frac{V_{ub}V_{us}^*}{V_{tb}V_{ts}^*} = -\left|\frac{V_{ub}V_{us}^*}{V_{tb}V_{ts}^*}\right| e^{i\alpha} = F_1 + iF_2\,, \tag{14}$$

where α is the unitarity triangle phase.

We also recall that the operator basis in $\mathcal{H}_{\text{eff}}$ is larger than what is shown in (2) in which the operators multiplying the CKM factor $V_{ub}V_{us}^*$ have been neglected. To calculate CP-asymmetry we have to put them back. Doing this, and using the unitarity relation $V_{cb}V_{cs}^* = -V_{ub}V_{us}^* - V_{tb}V_{ts}^*$, the effective Hamiltonian reads [28]

$$\mathcal{H}_{\text{eff}} = -\frac{G_F}{\sqrt{2}}\left\{\begin{array}{c} V_{tb}V_{ts}^*[C_7(\mu)O_7(\mu) + C_8(\mu)O_8(\mu) \\ +C_1(\mu)O_1(\mu) + C_2(\mu)O_2(\mu)] \\ \\ V_{ub}V_{us}^*[C_1(\mu)(O_{1u}(\mu) - O_1(\mu)) \\ +C_2(\mu)(O_{2u}(\mu) - O_2(\mu)) + \ldots] \end{array}\right\}. \tag{15}$$

In the above equation the ellipsis denotes the terms proportional to the Wilson coefficients $C_3 \ldots C_6$, and we have dropped them because they are very small as compared to C_1 and C_2. The operators O_{1u} and O_{2u} are defined as

$$O_{1u}(\mu) = \left(\bar{s}_L\gamma_\mu T^a u_L\right)\left(\bar{u}_L\gamma^\mu T_a b_L\right)\,,$$
$$O_{2u}(\mu) = \left(\bar{s}_L\gamma_\mu u_L\right)\left(\bar{u}_L\gamma^\mu b_L\right)\,.$$

The values of the Wilson coefficients in (15) are the same as we have already used in (8)–(10). Thus by including the annihilation contribution and also the effect of the operator O_{1u} and O_{2u}, the branching ratio from (6) can be written as

$$\begin{aligned}&\mathcal{B}_{\text{th}}\left(B^\pm \to K_1^\pm\gamma\right) = \tau_{B^+}\,\Gamma_{\text{th}}(B^\pm \to K_1^\pm\gamma) = \\ &\tau_{B^+}\frac{G_F^2\alpha|V_{tb}V_{ts}^*|^2}{32\pi^4}m_{b,\text{pole}}^2 M^3\left(1 - \frac{m_{K_1}^2}{M^2}\right)^3\left[\xi_\perp^{(K_1)}(0)\right]^2 \\ &\times\left\{\left(C_7^{(0)\text{eff}} + A_R^{(1)}\right)^2 + \left(F_1^2 + F_2^2\right)\left(A_R^u + L_R^u\right)^2 + 2F_1\right. \\ &\left.\times\left[C_7^{(0)\text{eff}}(A_R^u + L_R^u) + A_R^{(1)}L_R^u\right] \mp 2F_2\left[C_7^{(0)\text{eff}}A_I^u - A_I^{(1)}L_R^u\right]\right\},\end{aligned} \tag{16}$$

where $L_R^u = \epsilon_A C_7^{(0)\text{eff}}$ and the subscripts R and I denote the real and imaginary parts of the quantities involved. $A^{(1)}$ is the same as defined in (7), and A^u corresponds to the contribution from O_{1u} and O_{2u}, which can be written as

$$A^u(\mu) = \frac{\alpha_s(\mu)}{4\pi}C_2^{(0)}(\mu)\left[r_2(z) - r_2(0)\right] - \frac{\alpha_s(\mu_{\text{sp}})}{18\pi}C_2^{(0)} \times(\mu_{\text{sp}})\frac{\Delta F_\perp^{(K_1)}(\mu_{\text{sp}})}{\xi_\perp^{(K_1)}(0)}\frac{h^{(K_1)}(z,\mu_{\text{sp}})}{\langle\bar{u}^{-1}\rangle_\perp^{(K_1)}(\mu_{\text{sp}})}\,. \tag{17}$$

We now proceed to the numerical calculation of the branching ratios for the decay $B^+ \to K_1^+\gamma$. Using the value

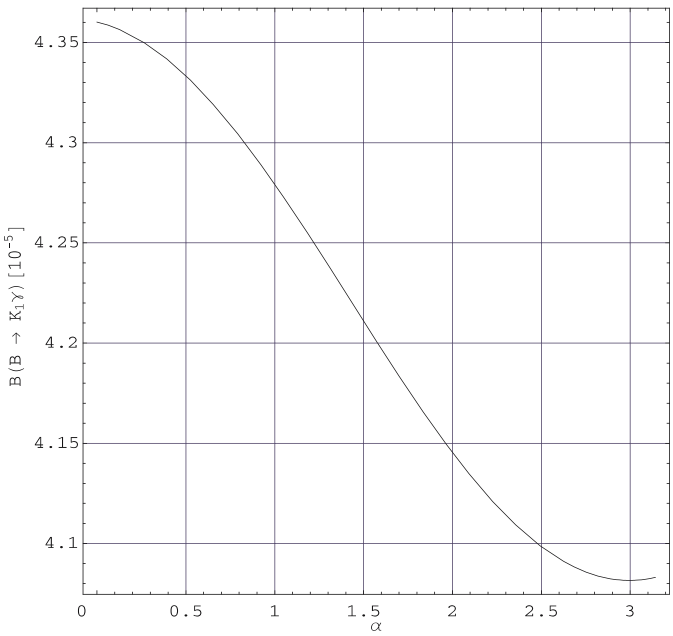

Fig. 7. Branching ratio for $B \to K_1\gamma$ decay versus the unitarity triangle phase α

of the CKM elements from [29], the values of $A^{(1)}(\mu)$ from [12] and the value of $C_2^{(0)}(\mu)$ from [13–15], the branching ratio is plotted with the unitarity triangle phase α as shown in Fig. 7.

One can easily see that varying the value of α in the range $77° \leq \alpha \leq 113°$ with $\alpha = 93°$ as the central value, there is a slight change in the value of the branching ratio for the decay $B \to K_1(1270)\gamma$ leaving the value of the form factor unchanged in this range as shown in the Fig. 8. We also note that the region of α where the branching ratio is effected is not allowed by the CKM unitarity constraints within the SM which typically yields $77° \leq \alpha \leq 113°$.

We now compute the leading order CP-asymmetry $\mathcal{A}_{\rm CP}\left(K_1^{\pm}\gamma\right)$ for the decay $B^{\pm} \to K_1^{\pm}\gamma$. The CP-asymmetry arises from the interference of the penguin operator O_7 and the four-quark operator O_2 [21–24]. The direct CP-asymmetry in the $B^{\pm} \to K_1^{\pm}\gamma$ is

$$\mathcal{A}_{\rm CP}\left(K_1^{\pm}\gamma\right) = \frac{\mathcal{B}\left(B^- \to K_1^-\gamma\right) - \mathcal{B}\left(B^+ \to K_1^+\gamma\right)}{\mathcal{B}\left(B^- \to K_1^-\gamma\right) + \mathcal{B}\left(B^+ \to K_1^+\gamma\right)} = \frac{2F_2\left(A_{\rm I}^u - \epsilon_A A_{\rm I}^{(1)}\right)}{C_7^{(0){\rm eff}}\left(1 + 2\epsilon_A\left[F_1 + \frac{1}{2}\epsilon_A\left(F_1^2 + F_2^2\right)\right]\right)} . \tag{18}$$

The dependence of the CP-asymmetry on the different parameters involved is shown in Figs. 9 and 10. In Fig. 9 we have plotted the CP-asymmetry versus the unitarity triangle phase α. It is seen that in the SM favored interval of α, $77° \leq \alpha \leq 113°$, the CP-asymmetry increases and reaches its maximum value which is 0.75% and is negative. This reduces to the value of 0.45% if one includes the hard-spectator corrections in addition to the vertex corrections and annihilation contributions.

Figure 10 shows the plot of $\mathcal{A}_{\rm CP}\left(K_1^{\pm}\gamma\right)$ with α at different values of the scale μ. It is very clear that the CP-asymmetry has a marked dependence on the scale μ. The value of the CP-asymmetry decreases from 0.8% to

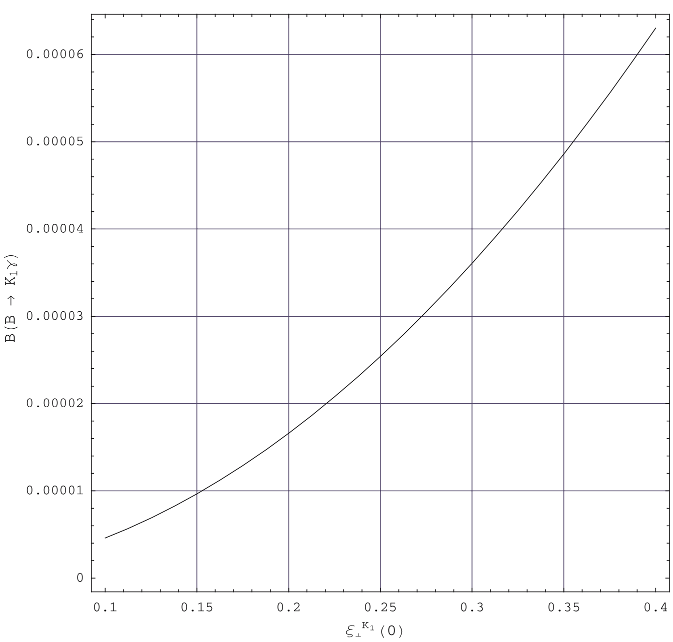

Fig. 8. Branching ratio for $B \to K_1\gamma$ decay versus the LEET form factor for the fixed value of $\alpha = 93°$

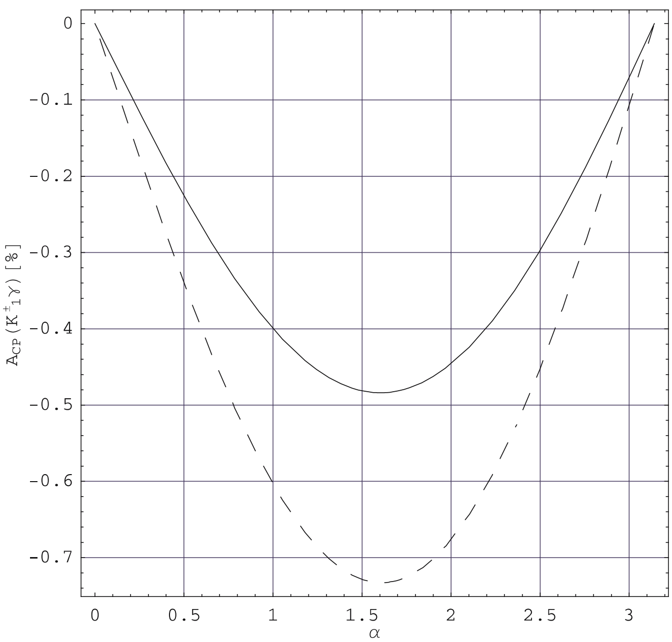

Fig. 9. CP-asymmetry ($\mathcal{A}_{\rm CP}\%$) versus the unitarity triangle phase α; the *dashed line* shows the value without hard-spectator correction and the *solid line* shows the value with hard-spectator correction

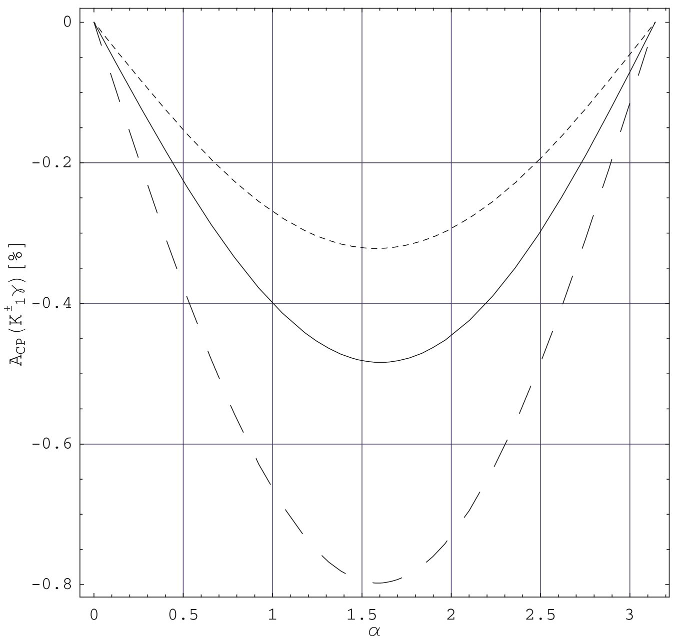

Fig. 10. CP-asymmetry ($\mathcal{A}_{\rm CP}\%$) versus the unitarity triangle phase α for different value of the scale μ; the *dashed line* shows the value at $m_{b,\rm pole}/2$; the *solid line* shows the value at $m_{b,\rm pole}$ and the *dotted line* shows it at $2m_{b,\rm pole}$

0.3% in the interval $m_{b,\rm pole}/2 \le \mu \le 2m_{b,\rm pole}$. A similar discussion for $B \to \rho\gamma$ is given in [10].

In conclusion, we have incorporated the effect of the annihilation and u-quark penguin contributions on the branching ratio for the decay $B \to K_1(1270)\gamma$. It is shown that the value of the LEET form factor remains the same even with inclusion of these annihilation contributions for the value of the unitarity triangle phase α favored by the standard model. Then, the CP-asymmetry $\mathcal{A}_{\rm CP}\left(K_1^{\pm}\gamma\right)$ for $B \to K_1(1270)\gamma$ has also been calculated. The CP-asymmetry received a contribution from the hard-spectator corrections which tend to decrease its value estimated from the vertex corrections alone. Unfortunately, the predicted value of the CP-asymmetry is sensitive to the choice of scale as well as to the quark mass ratio. The typical value of the CP-asymmetry lies around -0.5% which is almost same as for B to K^* decays.

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