

Implications of the first observation of $B \rightarrow K_1 \gamma$

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Implications of the recent new measurements of $B \rightarrow K_1 \gamma$ by Belle are examined. It is shown that the new branching ratio $\mathcal{B}[B \rightarrow K_1(1270)\gamma]$ requires very large form factor compared to the theoretically predicted one. This is an opposite case to $B \rightarrow K^* \gamma$ where theory expected larger branching ratio. Possible origins of the discrepancy are discussed.

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Radiative B decays to kaons provide a rich laboratory to test the standard model (SM) and probe new physics. $B \rightarrow K^* \gamma$ is a well established process among them. Higher resonant kaons such as $K_2^*(1430)$ are also measured by CLEO [1] and the B factories [2,3].

Recently, Belle has announced the first measurement of $K_1(1270)$ [4]:

$$\mathcal{B}[B^+ \rightarrow K_1^+(1270)\gamma] = (4.28 \pm 0.94 \pm 0.43) \times 10^{-5}. \quad (1)$$

There is also an upper bound on $K_1(1400)$ [4]:

$$\mathcal{B}[B^+ \rightarrow K_1^+(1400)\gamma] < 1.44 \times 10^{-5} \text{ (at 90\% C.L.)}. \quad (2)$$

There are many reasons to focus on the higher kaon resonances. First, they share lots of things with $B \rightarrow K^* \gamma$. At the quark level, both of them are governed by $b \rightarrow s \gamma$; all of the accumulated achievements of $b \rightarrow s \gamma$ can be used in radiative B decays to kaon resonances. For example, the same operators in the operator product expansion, the same corresponding Wilson coefficients, are available. In addition, when the hadronic descriptions are required, the resemblance between K^* and K_1 makes the analysis much easier. Especially, the light cone distribution amplitudes (DA) are same except the overall factor of γ_5 which gives rise to few differences in many calculations [5].

Second, $B \rightarrow K_{\text{res}}(\rightarrow K \pi \pi) \gamma$ can provide a direct measurement of the photon polarization [6]. In particular, it was shown that $B \rightarrow K_1(1400) \gamma$ can produce large polarization asymmetry of $\approx 33\%$ in the SM. In the presence of anomalous right-handed couplings, the polarization can be severely reduced in the parameter space allowed by current experimental bounds of $B \rightarrow X_s \gamma$ [7]. 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.

Third, theorists are now facing challenges from the discrepancy between their predictions and experiments. In fact, there have been noticeable theoretical advances in $B \rightarrow K^* \gamma$ over the last decade. QCD corrections at next-

to-leading order (NLO) of α_s was already considered in [8–10]. Furthermore, relevant Wilson coefficients have been improved [11,12] up to three-loop calculations. Recent developments of the QCD factorization (QCDF) [13] helped one calculate the hard spectator contributions systematically in a factorized form through the convolution at the heavy quark limit [14–16]. $B \rightarrow K^* \gamma$ is also analyzed in the effective theories at NLO, such as large energy effective theory [17] and the soft-collinear effective theory (SCET) [18].

But the nonperturbative analyses should be taken into account to complete the phenomenological explanation. QCD sum rule or the light cone sum rule (LCSR) is among the most reliable. It was pointed out in [17], however, that the LCSR results for the relevant form factor of $B \rightarrow K^* \gamma$ lead to a very large branching ratio compared to the measured one. Unfortunately, there is no way to explain the gap up to now.

The situation is more complicated in $B \rightarrow K_1 \gamma$. Based on the QCDF framework combined with the LCSR results, Ref. [5] predicted $\mathcal{B}[B^0 \rightarrow K_1^0(1270)\gamma] = (0.828 \pm 0.335) \times 10^{-5}$ and $\mathcal{B}[B^0 \rightarrow K_1^0(1400)\gamma] = (0.393 \pm 0.151) \times 10^{-5}$ at the NLO of α_s . New measurements (Eqs. (1) and (2)) certainly cast many questions about the theoretical predictions. Present work will be devoted to this issue.

The effective Hamiltonian for $b \rightarrow s \gamma$ is

$$\mathcal{H}_{\text{eff}}(b \rightarrow s \gamma) = -\frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* \sum_{i=1}^8 C_i(\mu) O_i(\mu), \quad (3)$$

where

$$\begin{aligned} O_2 &= (\bar{s}_i c_i)_{V-A} (\bar{c}_j b_j)_{V-A}, \\ O_7 &= \frac{em_b}{8\pi^2} \bar{s}_i \sigma^{\mu\nu} (1 + \gamma_5) b_j F_{\mu\nu}, \\ O_8 &= \frac{g_s m_b}{8\pi^2} \bar{s}_i \sigma^{\mu\nu} (1 + \gamma_5) T_{ij}^a b_j G_{\mu\nu}^a, \end{aligned} \quad (4)$$

are the relevant operators for present analysis. Here i, j are color indices, and we neglect the Cabibbo-Kobayashi-Maskawa (CKM) element $V_{ub} V_{us}^*$ as well as the s -quark mass. At next-to-leading order of α_s , the decay amplitude \mathcal{A} is given by

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$$\begin{aligned} \mathcal{A}(B \rightarrow K_1 \gamma) &= -\frac{G_F}{\sqrt{2}} V_{ib} V_{is}^* (C_7^{\text{eff}} \langle O_7 \rangle + C_2 \langle O_2 \rangle + C_8^{\text{eff}} \langle O_8 \rangle), \\ & \quad (5) \end{aligned}$$

where $\langle O_i \rangle \equiv \langle K_1 \gamma | O_i | B \rangle$. The leading contribution of $\langle O_7 \rangle$ is given by

$$\begin{aligned} \langle K_1(p', \epsilon) | \bar{s} i \sigma_{\mu\nu} q^\nu b | B(p) \rangle &= F_+^A(q^2) [(\epsilon^* \cdot q)(p + p')_\mu - \epsilon_\mu^* (p^2 - p'^2)] + F_-^A(q^2) [(\epsilon^* \cdot q) q_\mu - \epsilon_\mu^* q^2] \\ & \quad + \frac{F_0^A(q^2) \epsilon^* \cdot q}{m_B m} [(p^2 - p'^2) q_\mu - (p + p')_\mu q^2], \end{aligned} \quad (7a)$$

$$\langle K_1(p', \epsilon) | \bar{s} i \sigma_{\mu\nu} \gamma_5 q^\nu b | B(p) \rangle = i F_+^A(q^2) \epsilon_{\mu\nu\alpha\beta} \epsilon^{*\nu} q^\alpha (p + p')^\beta, \quad (7b)$$

where m and ϵ^μ are the mass and polarization vector of K_1 , respectively, and $q = p - p'$ is the photon momentum.

All the subleading contributions to $\langle O_7 \rangle$ shown in Fig. 1 are absorbed into the form factor F_+^A , while the corresponding Wilson coefficient C_7^{eff} contains its NLO parts,

$$C_7^{\text{eff}}(\mu) = C_7^{\text{eff}(0)}(\mu) + \frac{\alpha_s(\mu)}{4\pi} C_7^{\text{eff}(1)}(\mu). \quad (8)$$

On the other hand, the leading order $C_2^{(0)}$ and $C_8^{\text{eff}(0)}$ are sufficient for C_2 and C_8 since O_2 and O_8 contributions begin at NLO. The NLO contributions of $O_{2,8}$ can be written as

$$\langle O_i \rangle = \langle O_i \rangle_{VC} + \langle O_i \rangle_{HS} (i = 2, 8), \quad (9)$$

where $\langle O_i \rangle_{VC(HS)}$ are vertex corrections (hard spectator interactions) depicted in Figs. 2 (3). The branching ratio of $B \rightarrow K_1 \gamma$ is simply given by

$$\begin{aligned} \mathcal{B}(B \rightarrow K_1 \gamma) &= \tau_B \frac{G_F^2 \alpha m_b^2 m_B^3}{32\pi^4} \left(1 - \frac{m_A^2}{m_B^2}\right)^3 |F_+^A(0)|^2 |V_{ib} V_{is}^*|^2 \\ & \quad \times |C_7^{\text{eff}}(\mu_b) + A_{VC} + A_{HS}|^2. \end{aligned} \quad (10)$$

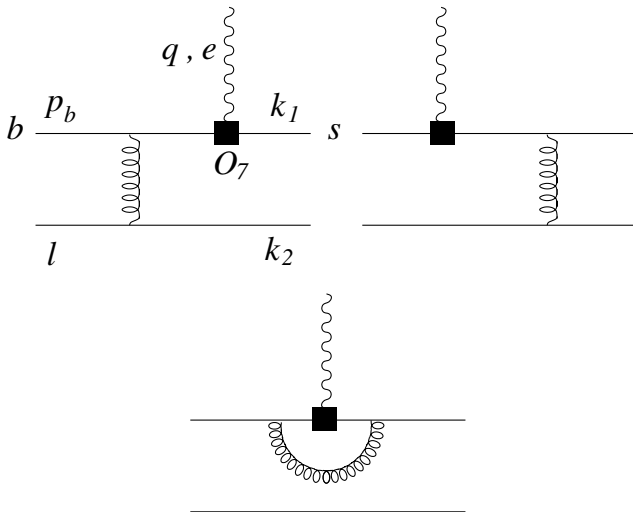


FIG. 1. NLO corrections to O_7 . These diagrams are absorbed into the weak form factor F_+^A .

$$\begin{aligned} \langle O_7 \rangle &\equiv \langle K_1(p', \epsilon) \gamma(q, e) | O_7 | B(p) \rangle \\ &= \frac{em_b}{4\pi^2} F_+^A(0) [\epsilon^* \cdot q (p + p') \cdot e^* - \epsilon^* \cdot e^* (p^2 - p'^2) \\ & \quad + i \epsilon_{\mu\nu\alpha\beta} \epsilon^{*\mu} \epsilon^{*\nu} q^\alpha (p + p')^\beta], \end{aligned} \quad (6)$$

with e^μ being the photon polarization vector. The form factor F_+^A is defined by

At the heavy quark limit,

$$\begin{aligned} A_{VC} &= \frac{\alpha_s(\mu_b)}{4\pi} \left\{ C_8^{\text{eff}}(\mu_b) \left[-\frac{32}{9} \ln \frac{m_b}{\mu_b} + \frac{4}{27} \right. \right. \\ & \quad \left. \left. \times (33 - 2\pi^2 + 6i\pi) \right] + C_2(\mu_b) \left[\frac{416}{81} \ln \frac{m_b}{\mu_b} + r_2 \right] \right\}, \\ A_{HS} &= \frac{4\pi\alpha_s(\mu_H) C_F}{N_c} \frac{f_B f_A^\perp}{\lambda_B m_B F_+^A(0)} \left\{ C_8^{\text{eff}}(\mu_H) \frac{1}{12} \langle u^{-1} \rangle_\perp \right. \\ & \quad \left. - C_2(\mu_H) \frac{1}{12} \left\langle \frac{\Delta i_5(z_0^{(c)}, 0, 0)}{\bar{u}} \right\rangle_\perp \right\}. \end{aligned} \quad (11)$$

See [5] for details.

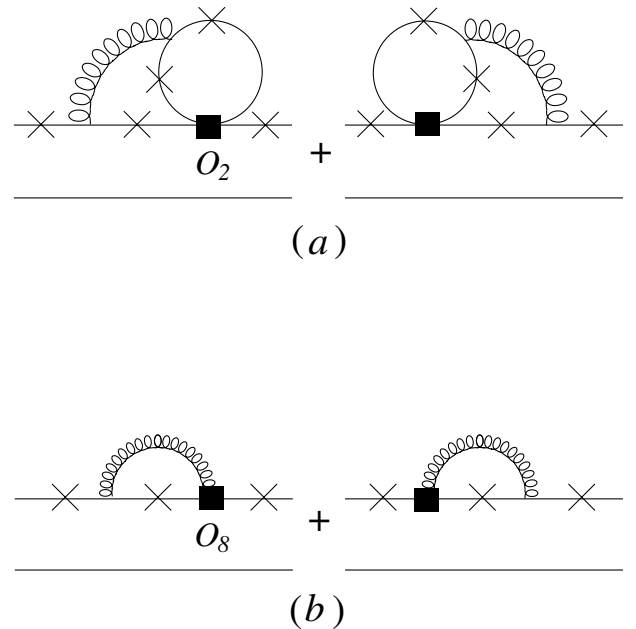


FIG. 2. Vertex corrections to the operators (a) O_2 and (b) O_8 . Crosses denote the possible attachment of the emitted photon.

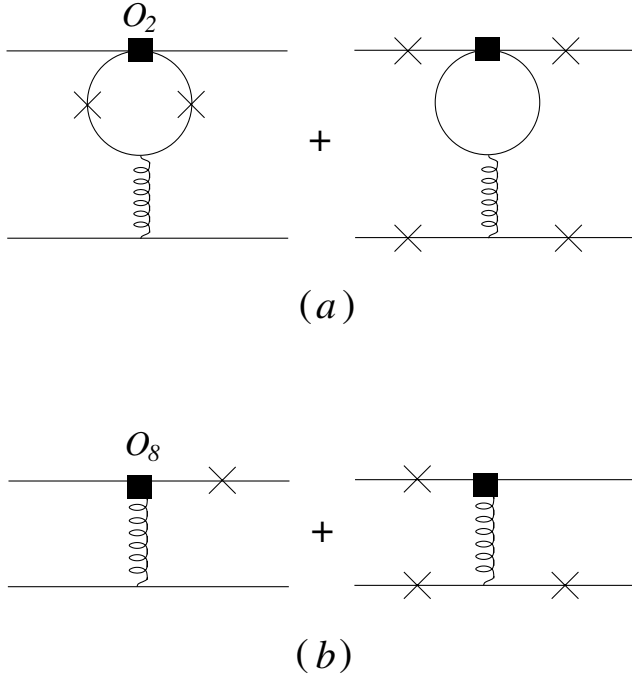


FIG. 3. Hard spectator interactions to (a) O_2 and (b) O_8 . First diagrams are leading contributions at the heavy quark limit.

Keeping the hadronic parameters specifically, we have

$$\begin{aligned} \mathcal{B}(B^0 \rightarrow K_1^0 \gamma) &= 0.003 \times \left(1 - \frac{m^2}{m_B^2}\right)^3 |F_+^A(0)| \\ &\times (-0.385 - i0.014) \\ &+ (f_A^1/\text{GeV})(-0.024 - i0.022)|^2, \end{aligned} \quad (12)$$

at the reference scales

$$\begin{aligned} (\mu_b, \mu_H) &= [m_b(m_b), \sqrt{\Lambda_H m_b(m_b)}] \\ &= (4.2 \text{ GeV}, 1.45 \text{ GeV}). \end{aligned} \quad (13)$$

It is now quite straightforward to extract the value of $F_+^A(0)$ from the new measurements (1) and (2). We have

$$F_+^{K_1(1270)}(0) = 0.32 \pm 0.03, \quad F_+^{K_1(1400)}(0) < 0.19, \quad (14)$$

where $f_{K_1(1270)} = 0.122 \text{ GeV}$, $f_{K_1(1400)} = 0.091 \text{ GeV}$ are used [19]. These must be compared with the LCSR results [19]

$$\begin{aligned} F_+^{K_1(1270)}(0)|_{\text{LCSR}} &= 0.14 \pm 0.03, \\ F_+^{K_1(1400)}(0)|_{\text{LCSR}} &= 0.098 \pm 0.02. \end{aligned} \quad (15)$$

Here we have another big difference between theory and experiment other than K^* . But the details of the differences

are quite opposite. In short,

$$F_{\text{theory}}^{K^*} > F_{\text{exp}}^{K^*}, \quad F_{\text{theory}}^{K_1} \ll F_{\text{exp}}^{K_1}. \quad (16)$$

There are some candidates to explain the discrepancy. Higher twist effects in the light cone DA are the first one. Usually they are process dependent, and are encoded in the coefficients of the Gegenbauer expansion. It is also known that they are asymptotically zero at $\mu \rightarrow \infty$ where μ is the renormalization scale. Ref. [17] estimated that the non-asymptotic correction of K^* at higher twist through the Gegenbauer moments to the operator O_8 is $\sim -20\%$. This is a bad news for $K_1(1270)$ if a similar tendency occurs for the axial Kaons since the present analysis is based on the asymptotic form of the light cone DA.

The second candidate is the nonzero mass effect. When calculating the hard spectator interactions in (11), it is assumed that the axial kaon is nearly massless and energetic. Although the assumption is acceptable for $m_{K_1} \ll m_B$, the mass hierarchy of $m_{K^*} < 1 \text{ GeV} < m_{K_1}$ might impose some doubts about the common framework for both K^* and K_1 . Note that the chiral symmetry is broken around 1 GeV.

But including nonzero mass corrections is very non-trivial. Since the relevant large scale in $B \rightarrow K_1 \gamma$ is m_B , possible mass corrections will appear in the form of m_{K_1}/m_B . It means that to fully appreciate the mass effects, one has to consider the $1/m_B$ (or $1/m_b$) corrections throughout the analysis, which is not well established so far. Since the discrepancy of (16) is quite large and $m_{K_1}/m_B \approx 0.24$, one should expect large corrections like chiral enhancement in nonleptonic decays at $1/m_B$.

Third, the framework of QCDF might not be adequate for the axial kaons. The main idea of QCDF can be summarized by [16]

$$\begin{aligned} \langle V \gamma(\epsilon) | O_i | B \rangle &= \left[F^{B \rightarrow V}(0) T_i^I \right. \\ &\left. + \int_0^1 d\xi dv \Phi_V(v) T_i^{II}(\xi, v) \Phi_B(\xi) \right] \cdot \epsilon, \end{aligned} \quad (17)$$

where $T_i^{I,II}$ are the hard scattering kernels. The kernel T_i^{II} is concerned with the hard spectator interactions. The factorization of (17) holds when the hard kernels are perturbatively calculable. All the nonperturbative physics are encapsulated in the DAs. A great discrepancy of (16) suggests that this may not be the case for K_1 . Some of the simple model calculations based on the heavy quark effective theory (HQET) predict rather large branching ratios; see the Table V in [5] or refer to [20]. Since the higher resonant kaons are heavy $\geq 1 \text{ GeV}$, it is quite natural and attractive to consider them as heavy mesons. In the heavy quark scheme, hard spectator interaction is

inconceivable since almost all the momentum of initial heavy quark is transferred to the final one.

We can also question the reliability of the QCD sum rule or LCSR results. It is a common knowledge that the stability of an observable against the Borel parameter in the QCD sum rule gets poorer as higher resonances are involved. Still, the problem of how to describe the higher kaon resonances remains. It is also noticeable that the lattice calculation is very close to the QCD sum rule result for K^* [17,21]. Much more reliable nonperturbative analyses are required in the near future.

Next, possible mixing in $K_1(1270)$ and $K_1(1400)$ cannot explain the large mismatches of (16). Quark model states 3P_1 and 1P_1 can mix to form physical states $K_1(1270)$ and $K_1(1400)$. The form factors are now written as [22]

$$\begin{aligned} F_+^{K_1(1270)}(0) &= Y_A(0) \sin\theta + Y_B(0) \cos\theta, \\ F_+^{K_1(1400)}(0) &= Y_A(0) \cos\theta - Y_B(0) \sin\theta, \end{aligned} \quad (18)$$

where $Y_{A,B}$ are the form factors corresponding to the angular momentum eigenstates. The enhancement from maximal mixing is only a factor of $\sqrt{2}$, assuming $Y_A(0) \approx$

$Y_B(0)$. A substantial growth in $Y_{A,B}(0)$ is inevitable to explain the experimental data. On the other hand, the usefulness of mixing lies in the fact that it can naturally explain a strong suppression of $B \rightarrow K_1(1400)\gamma$. But it is too early to say something about this point with the new upper bound of (2); the LCSR result (15) is still within the boundary. Therefore, a new observation of $B \rightarrow K_1(1400)\gamma$ is much anticipated.

Finally, it is quite unlikely that the annihilation topology would give considerable contributions, as pointed out in [15,16].

In conclusion, we surveyed the implications of the first observation of $B \rightarrow K_1\gamma$. The values of the relevant form factors are extracted from the experimental data at NLO of α_s . We found that a very large discrepancy between theory and experiment is reproduced after $B \rightarrow K^*\gamma$. Eliminating the gap will be a great challenge in theory. Further observation by other B factory as well as of $K_1(1400)$ will provide much interest in coming days.

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- [1] CLEO Collaboration, T. E. Coan *et al.*, Phys. Rev. Lett. **84**, 5283 (2000).
 - [2] Belle Collaboration, S. Nishida *et al.*, Phys. Rev. Lett. **89**, 231801 (2002).
 - [3] BABAR Collaboration, B. Aubert *et al.*, Phys. Rev. D **70**, 091105 (2004).
 - [4] Belle Collaboration, K. Abe *et al.*, hep-ex/0408138.
 - [5] J. P. Lee, Phys. Rev. D **69**, 114007 (2004).
 - [6] M. Gronau, Y. Grossman, D. Pirjol, and A. Ryd, Phys. Rev. Lett. **88**, 051802 (2002); M. Gronau and D. Pirjol, Phys. Rev. D **66**, 054008 (2002).
 - [7] J. P. Lee, Phys. Rev. D **69**, 014017 (2004); hep-ph/0312010.
 - [8] J. M. Soares, Nucl. Phys. B **367**, 575 (1991); Phys. Rev. D **49**, 283 (1994).
 - [9] C. Greub, H. Simma, and D. Wyler, Nucl. Phys. B **434**, 39 (1995); **444**, 447(E) (1995).
 - [10] C. Greub, T. Hurth, and D. Wyler, Phys. Rev. D **54**, 3350 (1996).
 - [11] K. Adel and Y. P. Yao, Phys. Rev. D **49**, 4945 (1994).
 - [12] K. G. Chetyrkin, M. Misiak, and M. Munz, Phys. Lett. B **400**, 206 (1997); **425**, 414(E) (1998).
 - [13] M. Beneke, G. Buchalla, M. Neubert, and C. T. Sachrajda, Nucl. Phys. B **591**, 313 (2000).
 - [14] M. Beneke and T. Feldmann, Nucl. Phys. B **592**, 3 (2001).
 - [15] M. Beneke, T. Feldmann, and D. Seidel, Nucl. Phys. B **612**, 25 (2001).
 - [16] S. W. Bosch and G. Buchalla, Nucl. Phys. B **621**, 459 (2002).
 - [17] A. Ali and A. Y. Parkhomenko, Eur. Phys. J. C **23**, 89 (2002).
 - [18] J. Chay and C. Kim, Phys. Rev. D **68**, 034013 (2003).
 - [19] A. S. Safir, Eur. Phys. J. C **3**, 15 (2001).
 - [20] A. Ali, T. Ohl, and T. Mannel, Phys. Lett. B **298**, 195 (1993).
 - [21] UKQCD Collaboration, L. Del Debbio, J. M. Flynn, L. Lellouch, and J. Nieves, Phys. Lett. B **416**, 392 (1998).
 - [22] H. Y. Cheng and C. K. Chua, Phys. Rev. D **69**, 094007 (2004).