Quark off-shell effects in $B \rightarrow \gamma\gamma$ decays

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Abstract

We demonstrate that in the decay of a heavy-light quark system, like the process $B_s \rightarrow \gamma\gamma$ under consideration, there is an important contribution coming from operators that vanish on the free quark mass-shell. Although numerical estimates are model-dependent, the significance of genuine off-shell effects is quite general, and independent of the details of the bound-state dynamics. Such contributions will typically be somewhat smaller, but still significant in magnitude compared with the contribution from the standard off-diagonal, $b \rightarrow s\gamma$ magnetic-moment operator.

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1 Introduction

With the advent of heavy-quark effective theories [1] more study has been devoted also to radiative heavy-light quark system decays, such as $B_s \rightarrow K^* \gamma$ [2, 3, 4]. Radiative flavour-changing decays thereby appear to be ideally suited for learning some details of the corresponding meson-quark dynamics. For example, the study of the radiative $b \rightarrow s + \gamma$ decay [5] may reveal the effect of the Fermi motion of the heavy quark. This process also raises rather high expectations in constraining various extensions beyond the standard model [6]. In the process $b \rightarrow s + \gamma$ and its hadronic $B_s \rightarrow K^* \gamma$ counterpart, the light quark is merely a spectator. However, additional effects from the motion of the light quark might be expected in flavour-changing two-photon decays, such as $b \rightarrow s \gamma \gamma$ and its hadronic $B_s \rightarrow \gamma \gamma$ version, where the light and the heavy quark fuse together. Operators for the electroweak flavour change in $q \rightarrow q' \gamma \gamma$ transitions contain a genuine off-shell piece, which vanishes on the free-quark mass shell [7, 8]. It is clear that such an on-shell limit is inapplicable to quarks immersed in light mesons, such as K^0 . Consequently, we first revealed [9] the significant off-shell contribution to the CP-violating $K_L \rightarrow \gamma \gamma$ amplitude and, more recently, to the CP-conserving $K_L \rightarrow \gamma \gamma$ amplitude [10]. The theoretical framework which enabled us to treat off-shell quarks was the chiral quark model [11]. However, the essential features of the off-shell appearance are independent of the particular model of the meson-quark interaction, as demonstrated explicitly in the bound-state approach [12].

In the present letter we want to go one step further, by demonstrating how the same type of off-shellness appears in heavy-light systems, such as the B_s meson, by considering $B_s \rightarrow 2\gamma$. One might expect such genuine off-shell effects to be suppressed by $\bar{\Lambda}/m_b$, where $\bar{\Lambda} = M_B - m_b$ is the mass difference between the physical B -meson and the mass of the current b -quark. By a convenient rearranging of the off-shell radiative transition operator we demonstrate the departure from what would correspond to the mutual 1PI and 1PR diagram cancellation in the free-quark limit [8]. There is a subtle interplay of the genuine 1PI local two-photon flavour-changing transition and its one-photon flavour-changing counterpart when they are embedded in the constituent/bound-quark loop. A piece of the latter, which cancels the local two-photon flavour-changing transition, is extracted. Here we present the remaining, uncanceled part of the 1PR amplitude. This remaining contribution contains two pieces: the conventional off-diagonal magnetic-moment transition operator and the genuinely off-shell part.

After presenting the aforementioned facts in detail, we turn to their explicit realization in evaluating the $B_s \rightarrow 2\gamma$ decay amplitude.

2 Off-shellness in the radiative $b \rightarrow s$ transition

In our previous papers we proposed a simple parametrization of the off-shell effects accompanying flavour-changing radiative transitions. The one-loop level electroweak transitions $b \rightarrow s\gamma$ and $b \rightarrow s\gamma\gamma$, related by the Ward identities, can be comprised as

$$\mathcal{L}(b \rightarrow s)_\gamma = B(M_W) \epsilon^{\mu\nu\lambda\rho} F_{\mu\nu} (\bar{s}_L i \overleftrightarrow{D}_\lambda \gamma_\rho b_L), \quad (1)$$

where the coefficient $B(M_W) \sim eG_F\lambda_{KM}$ contains the result of the relevant loop diagrams, and λ_{KM} is the relevant KM factor for $b \rightarrow s$ transitions. The covariant derivative contains both the gluon and photon fields. One can decompose this expression into two pieces, $\mathcal{L}(s \rightarrow d)_\gamma = \mathcal{L}_F + \mathcal{L}_\sigma$. The term

$$\mathcal{L}_F = B_F \bar{s} [(i\gamma \cdot D - m_s) \sigma_{\mu\nu} F^{\mu\nu} L + \sigma_{\mu\nu} F^{\mu\nu} R (i\gamma \cdot D - m_b)] b \quad (2)$$

would vanish by applying the (perturbative) equation of motion (or in momentum space, for on-shell free quarks), i.e for $(i\gamma \cdot D - m_{b,s}) \rightarrow 0$. (In (2) it is understood that the covariant derivative acts on the nearest fermion field to the left or right.) In this case $\mathcal{L}(b \rightarrow s)_\gamma$ reduces to its on-shell remnant

$$\mathcal{L}_\sigma = B_\sigma \bar{s} (m_b \sigma_{\mu\nu} F^{\mu\nu} R + m_s \sigma_{\mu\nu} F^{\mu\nu} L) b, \quad (3)$$

the familiar off-diagonal magnetic moment term. This conforms to previous results on flavour-changing two-photon transitions [13] that only the off-diagonal magnetic-moment term in convolution with another ordinary photon emission contributes in the partonic (free-quark mass-shell) picture. Oppositely, the departure from the quark mass-shell regime can be measured by the $\mathcal{L}_F$ term, containing also the direct 1PI flavour-changing two-photon transition. The natural environment for such a study of the off-shellness is flavour transitions within mesons, which requires to scale the operator $\mathcal{L}_F$ defined at the M_W scale, down to the B -meson scale. The expressions (2) and (3) take into account that the coefficients, B_F of $\mathcal{L}_F$ and B_σ of $\mathcal{L}_\sigma$, both being equal to B at the W -scale, evolve differently down to the $\mu = m_b$ scale. This difference between B_F and B_σ is due to different anomalous dimensions of the operators in (2) and (3). By repeating the procedure of ref. [14], giving for the off-diagonal magnetic moment operator

$$B_\sigma(\mu) = \eta^{-16/3} \left[B(M_W) + \frac{58}{135} (\eta^{10/3} - 1) + \frac{29}{189} (\eta^{28/3} - 1) \right], \quad (4)$$

we obtain for the off-shell operator (having zero anomalous dimension):

$$B_F(\mu) = B(M_W) - \frac{58}{81} (\eta^{-6/23} - 1) + \frac{29}{81} (\eta^{12/23} - 1), \quad (5)$$

where $\eta = \alpha_s(\mu)/\alpha_s(M_W)$. This result, obtained within a simplified calculation (with a truncated basis) is in agreement with the QCD corrections found in [15].

By employing the value of the quantity η at the b -quark scale, we obtain

$$\frac{B_F}{B_\sigma} \simeq 4/3 \quad (\mu = m_b) . \quad (6)$$

Bound-state models often involve the constituent quark mass instead of the current one. Therefore, we make another decomposition of the total radiative-decay Lagrangian

$$\mathcal{L}(b \rightarrow s)_\gamma = \tilde{\mathcal{L}}_F + \tilde{\mathcal{L}}_\sigma , \quad (7)$$

where $\tilde{\mathcal{L}}_F$ is obtained from $\mathcal{L}_F$ simply by replacing the current masses (m_q) by the constituent ones (M_q):

$$\tilde{\mathcal{L}}_F = \mathcal{L}_F(m_q \rightarrow M_q) , \quad (8)$$

and $\tilde{\mathcal{L}}_\sigma$ is the remainder

$$\tilde{\mathcal{L}}_\sigma = \bar{s} (\tilde{B}_R \sigma_{\mu\nu} F^{\mu\nu} R + \tilde{B}_L \sigma_{\mu\nu} F^{\mu\nu} L) b , \quad (9)$$

where

$$\tilde{B}_L = B_F(M_s - m_s) + B_\sigma m_s ; \quad \tilde{B}_R = B_F(M_b - m_b) + B_\sigma m_b . \quad (10)$$

Contrary to the current quark masses $m_{b,s}$ sitting in the quark propagators in the asymptotically free regime, the masses $M_{b,s}$ sit in the poles of the quarks propagating in the B -meson.

To calculate the $B \rightarrow 2\gamma$ amplitude from the Lagrangian $\mathcal{L}(b \rightarrow s)_\gamma$ one needs some bound-state model. Let us in this section be quite general, and just assume that such a meson-quark interaction exists. Moreover, we find it suitable to split the generically off-shell operator $\tilde{\mathcal{L}}_F$ into the one- and two-photon pieces, $\tilde{\mathcal{L}}_F = \tilde{\mathcal{L}}_F^{1\gamma} + \tilde{\mathcal{L}}_F^{2\gamma}$, where

$$\tilde{\mathcal{L}}_F^{2\gamma} = e_D B_F \bar{s} [A \cdot \gamma \sigma^{\mu\nu} F_{\mu\nu} L + \sigma^{\mu\nu} F_{\mu\nu} R \gamma \cdot A] b , \quad (11)$$

and the one photon piece is

$$\tilde{\mathcal{L}}_F^{1\gamma} = B_F \bar{s} [(i\partial \cdot \gamma - M_s) \sigma^{\mu\nu} F_{\mu\nu} L + \sigma^{\mu\nu} F_{\mu\nu} R (i\gamma \cdot \partial - M_b)] b . \quad (12)$$

The piece $\tilde{\mathcal{L}}_F^{2\gamma}$ gives rise to the first quark loop diagram in Fig. 1. The corresponding amplitude $A^{2\gamma}$ can be split in two parts corresponding to the L - and R - part of $\tilde{\mathcal{L}}_F^{2\gamma}$:

$$A^{2\gamma} = A_L^{2\gamma} + A_R^{2\gamma} . \quad (13)$$

The amplitude due to $\tilde{\mathcal{L}}_F^{1\gamma}$ is given by the last two diagrams in Fig. 1, the two respective pieces in

$$A^{1\gamma} = A_{L+R}^s + A_{L+R}^b. \quad (14)$$

A_{L+R}^q refers to the diagram where the extra photon (not contained in $\tilde{\mathcal{L}}_F^{1\gamma}$) is emitted by a $q = s, b$ quark. Then, using $SS^{-1} = S^{-1}S = 1$ for the inverse propagator $(\gamma \cdot p_q - M_q) = S_q^{-1}(p_q)$ in $\tilde{\mathcal{L}}_F^{1\gamma}$, we obtain

$$A_{L+R}^s = -A_L^{2\gamma} + \bar{A}_R^s; \quad A_{L+R}^b = -A_R^{2\gamma} + \bar{A}_L^b. \quad (15)$$

Thus the contribution due to $\tilde{\mathcal{L}}_F^{2\gamma}$ in (13) cancels with a part of the contributions in (14), coming from $\tilde{\mathcal{L}}_F^{1\gamma}$, and the total amplitude due to $\tilde{\mathcal{L}}_F$ is given by

$$A_{Tot}^F = A^{1\gamma} + A^{2\gamma} = \bar{A}_L^b + \bar{A}_R^s. \quad (16)$$

The reduced amplitudes $\bar{A}_L^b$ and $\bar{A}_R^s$ correspond to the diagrams in Fig. 2, involving effective $B_s b\bar{b}\gamma$ and $B_s s\bar{s}\gamma$ vertices. This result is equivalent to that of ref. [10], where it was shown that if $\mathcal{L}_F$ was transformed away into the wave function, it will reappear in the bound state low-energy sector. That is, the diagrams of Fig. 2 are the part of the off-shell operator that cannot be transformed away. This result is *completely general*, and does *not depend on the details of the bound state dynamics*, i.e. the form of G_B . Thus, it should be emphasized that although there is a cancellation between $\tilde{\mathcal{L}}_F^{2\gamma}$ and a part of $\tilde{\mathcal{L}}_F^{1\gamma}$, there is a remaining non-zero genuine off-shell contribution which has been neglected in the literature by appealing to the (perturbative) equations of motion.

Let us also note that the loop integrals in $\bar{A}_L^b$ and $\bar{A}_R^s$ are multiplied by M_s and M_b , respectively. Then we see from the structure of $\tilde{\mathcal{L}}_F$ that the off-shellness of the b - and s -quarks in the B_s meson are of order M_s and M_b , respectively: The off-shellness of the light quark is characterized by the heavy-quark mass, whereas the off-shellness of the heavy quark is characterized by the light-quark mass. For the s - and b -quarks in play,

$$(\gamma \cdot p_b - M_b) \sim M_s \quad \text{and} \quad (\gamma \cdot p_s - M_s) \sim M_b, \quad (17)$$

where p_q are the quark-momenta for $q = b, s$. This is an interesting and plausible result. However, as discussed in the next section, there is a competing effect from bound-state dynamics, namely a damping of the light-quark momenta [16]. This means that the loop integral multiplying M_b is damped compared to the one multiplying M_s , and the effect of the off-shellness of the light s -quark in $\mathcal{L}_F$ is less pronounced.

In a recent paper [4], it has been pointed out that there are off-shell bound-state effects in the process $B \rightarrow K^* \gamma$. However, only the off-shellness

of the b -quark was taken into account, whereas we find that for $B_s \rightarrow 2\gamma$ the off-shell effect of the s -quark cannot be neglected.

In the next section we will elaborate the $B_s \rightarrow 2\gamma$ amplitude to some detail.

3 The quark-loop $B_s \rightarrow \gamma\gamma$ amplitude

As for $K \rightarrow \gamma\gamma$ amplitudes [9, 10], the flavour-changing radiative vertices have to be supplemented by a quark-meson interaction in order to perform the full quark-loop evaluation. In contradistinction to the analogous $K \rightarrow \gamma\gamma$ decay, the heavy B -meson cannot be treated as a Goldstone boson of the chiral quark model adopted earlier. However, the pseudoscalar character of B -mesons allows us to parametrize this interaction in a simple way, replacing the $\mathcal{L}_\chi$ term of ref. [10] by

$$iG_B \bar{s} \gamma_5 b B_s . \quad (18)$$

This interaction may in general be non-local, i.e. G_B might be momentum dependent [16]. Thereby, as usually done [9], we trade the meson-quark coupling G_B in favour of the meson-decay constant f_B . In calculating the contributions from $\mathcal{L}_F$ and $\mathcal{L}_\sigma$ in (2) and (3) respectively, we obtain an amplitude of the following form :

$$M(B_s \rightarrow \gamma\gamma) = e_D f_B [A_{(+)} F_{\mu\nu} F^{\mu\nu} + i A_{(-)} F_{\mu\nu} \tilde{F}^{\mu\nu}] , \quad (19)$$

$$A_{(\pm)} = \tau_F^{(\pm)} B_F + \tau_\sigma^{(\pm)} B_\sigma , \quad (20)$$

where the quantities $\tau_{F,\sigma}^{(\pm)}$ are dimensionless and depend on the bound-state dynamics. They turn out to be of order one. The coefficients B_F of $\mathcal{L}_F$ and B_σ of $\mathcal{L}_\sigma$, now contain the KM factors relevant for the $b \rightarrow s$ transition and are renormalized at the scale $\mu = m_b$.

The quantities F and $\tilde{F}$ are the photonic field strength and its dual. Here, in calculating the decay rate, the contributions are collected in two characteristic final-state pieces: the $F \cdot F$ part and the $F \cdot \tilde{F}$ part. Note that the $F \cdot F$ part was not present in the previously calculated direct $\mathcal{O}(p^2)$ K_L -decay amplitude [9, 10]. In the formal limit where the current quark masses $m_{b,s} \rightarrow 0$, the $F_{\mu\nu} \tilde{F}^{\mu\nu}$ term of (19) should reduce to the anomalous $\bar{K}^0 \rightarrow 2\gamma$ amplitude. However, in the real world $M_b \gg M_s \gtrsim M \sim 300$ MeV, and the result for $\text{Br}(B_s \rightarrow 2\gamma)$ will be rather different from $\bar{K}^0 \rightarrow 2\gamma$.

Because we want, in this paper, to demonstrate the significance of the genuine off-shell contribution, and do not aim at an exact quantitative treatment, we have for illustrative purposes considered two different cases to estimate the quantities $\tau_{F,\sigma}^{(\pm)}$: 1) The limit of a constant G_B (local interaction), and 2) a form-factor damping of the light quark momenta as in [16].

1) Using the approximation $M_b \gg M_s$, we obtain:

$$f_B = N_c \frac{G_B}{4\pi^2} \frac{M_b}{2} \hat{f}, \quad (21)$$

where the $\hat{f}$ contains the loop-evaluation result which, in dimensional regularization is

$$\hat{f} = \frac{1}{\varepsilon} - \gamma_E - \ln \left(\frac{M_b^2}{4\pi\mu^2} \right) + 1. \quad (22)$$

Note that in the $K \rightarrow 2\gamma$ we used an expansion in the K -meson momenta, corresponding to a chiral expansion. In the B -meson case, we cannot do an expansion in the meson momentum k because $k^2 = M_B^2 \simeq M_b^2$ is the big scale in the matrix element.

In the extreme limit where M_s is neglected and we only keep the leading terms that can be incorporated in f_B , we obtain

$$\tau_\sigma^{(\pm)} = \tau_F^{(+)} = -\tau_F^{(-)} = 1, \quad (23)$$

giving a branching ratio

$$\text{Br}(B_s \rightarrow 2\gamma) \simeq 2 \times 10^{-8},$$

slightly smaller than the free-quark estimate [18], which is proportional to the inverse of the light s -quark mass. To be a little more detailed, the contribution from $\mathcal{L}_\sigma$ is not directly proportional to f_B in (21), (22), but to

$$\hat{f} + \eta + \pi + \frac{1}{2} \ln(M_b^2/M_s^2).$$

where $\eta = \frac{3}{2}$ in the 't Hooft-Veltman scheme and $\eta = \frac{1}{2}$ in naive dimensional regularization. The contribution from $\mathcal{L}_F$ is proportional to $(\hat{f} - 1)$. Thus it seems that the σ term is more important than the F term before QCD corrections are taken into account.

2) Using $G_B = H_B (\Lambda_s^2 - p_s^2)^{-1}$, we find (for the rather small value $\Lambda_s = M_s$, which we use for simplicity):

$$f_B \simeq N_c \frac{H_B}{4\pi^2} \frac{1}{2M_b} \left(\ln \frac{M_b^2}{M_s^2} + 1 \right). \quad (24)$$

This gives for $\mathcal{L}_\sigma$

$$\tau_\sigma^{(\pm)} = \tau_\sigma \simeq \frac{M_b^2}{4M_s^2} \left(\ln \frac{M_b^2}{M_s^2} + 1 \right)^{-1}. \quad (25)$$

For $\mathcal{L}_F$ we obtain

$$\tau_F^{(\pm)} = 2 \frac{M_s}{M_b} \left(1 \pm \frac{M_s}{M_b} \ln \frac{M_b^2}{M_s^2} \right) \tau_\sigma + \tau_{F,An}^{(\pm)}, \quad (26)$$

where $\tau_{F,An}^{(+)} = 0$ by definition, and from [9] we get the anomalous term $\tau_{F,An}^{(-)} \gtrsim 1$ obtained in dimensional regularization, which cannot be described properly in a form factor scheme[16]. In case 2) the leading term will be inversely proportional to the square of the constituent s -quark mass M_s . Qualitatively, this result is not far from that given in [18]. However, the small value $\Lambda_s = M_s$, chosen for illustrative purposes, overestimates τ_σ . For larger and more realistic values of Λ_s , we obtain $\tau_\sigma \sim 2$ to 3. We find that $\text{Br}(B_s \rightarrow 2\gamma)$ is increased in case 2) with respect to case 1). In general, we find that the genuine off-shell term $\mathcal{L}_F$ increases the rate by a factor $\sim$ of 1.5 to 3. We conclude that values of the order 10^{-8} and 10^{-7} are realistic. Our (rough) prediction is still two orders of magnitude above the LD estimates based on the vector-meson dominance [19].

We observe that in case 2) above, the genuine off-shell contribution is suppressed by $1/M_b$ in accordance with the expectations mentioned in the introduction. In [4] the genuine off-shell amplitude has a phase that almost washes out the effect in the branching ratio. Such possible phases have not been taken into account here [16]. It should be emphasized that even if more sophisticated and more accurate calculations in terms of bound states [20, 12] or sum rules [3] could have been performed, the simple considerations above serves our purpose, namely to demonstrate the significance of $\mathcal{L}_F$.

4 Conclusions

We have demonstrated the significance of quark off-shell effects in flavour-changing two-photon decays $b \rightarrow s\gamma\gamma$ and its hadronic $B_s \rightarrow \gamma\gamma$ counterpart. Thus, the same basic off-shell effect previously shown to take place in $\bar{K}^0 \rightarrow \gamma\gamma$ also takes place in a heavy-light bound-state calculation. There is a 2γ electroweak 1PI transition, which cancels with a part of the 1PR transition, but some effect of $\mathcal{L}_F$ remains.

The remaining contribution from the off-shell operator $\mathcal{L}_F$ corresponds to loop diagrams containing effective $B_s \bar{b}b\gamma$ and $B_s \bar{s}s\gamma$ vertices. This result is equivalent to the one presented in [9]: $\mathcal{L}_F$ may be transformed away into the wave function, but it then reappears in the bound-state dynamics and gives a physical effect.

Within our rather qualitative approach, we find a substantial contribution to $B_s \rightarrow \gamma\gamma$ from the genuine off-shell piece $\mathcal{L}_F$ of the effective Lagrangian. Although the hadronic matrix element is model dependent, there are substantial off-shell contributions, which increase the rate by a factor of ~ 1.5 to 3, giving a branching ratio of order 10^{-8} to 10^{-7} . Maybe a calculation of $\mathcal{L}_F$ within some variant of sum rules [3] could resolve some of the numerical uncertainties.

Our qualitative conclusion is clear: The genuine off-shell effect is substantial, and both the off-shellness of the heavy b -quark and the light s -quark should be considered.

More generally, we have demonstrated that the same basic off-shell effect takes place in different flavour-changing two-photon decays, such as $\bar{K}^0 \rightarrow \gamma\gamma$ and $B_s \rightarrow \gamma\gamma$, although they belong to such different environments as the chiral perturbation theory and heavy-light bound-state calculation.

We have demonstrated that the naive use of the (perturbative) equations of motion are not applicable in general. This should be no surprise, because the low-energy interactions, represented by (18) and the photon field sitting in the covariant derivative, might be interpreted as external fields. Thus the (perturbative) equations of motion are changed.

The genuine off-shell effects are formally suppressed in a certain limit by $1/M_b$ for $B \rightarrow \gamma\gamma$. However, the suppression is not very pronounced numerically.

In view of the attempts to constrain the various possible extensions of the standard model from the studies of the $B_s \rightarrow K^*\gamma$ decay, one should be aware of the off-shell effects described in the present paper.

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Figure captions:

Fig 1. : The two-photon contribution and the one-photon contributions to $B_s \rightarrow 2\gamma$. The cross within the circle denotes the effective vertex corresponding to $\tilde{\mathcal{L}}_F$.

Fig. 2: The remaining contribution to $B_s \rightarrow 2\gamma$ after the cancellation of the irreducible two-photon quark diagrams.

Fig. 1

Fig. 2