Soft gluon contributions to the $B \to K\eta'$ amplitude in a low energy bosonization model

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July 15, 2017

Abstract

Intriguing $B \to K\eta'$ decays provide a unique opportunity to study a joining of two-gluon configurations arising from the penguin $b \to sG$ and $b \to sGG$ transitions, with those inherent to the $\eta'$ particle. We employ the heavy-light chiral quark model, applied previously to a somewhat related $B \to D\eta'$ decay, as a calculational tool accounting for the nonperturbative soft gluon contributions to the amplitude at hand. Thereby we arrive at a novel contribution to the singlet penguin amplitude, which within our model accounts for $\sim 10\%$ of the measured $B \to K\eta'$ amplitude.

PACS: 12.15.Ji; 12.39.-x; 12.39.Fe; 12.39.Hg
Keywords: B mesons, Rare decays, Heavy quarks, Chiral Lagrangians

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1 Introduction

The data on rare decays from CLEO, BaBar, Belle and Tevatron hold promise for deepening our understanding of the interplay between flavour-changing and QCD dynamics. In particular, the surprise regarding the $B \to K \eta'$ amplitude, existing already for considerable time [1, 2], calls for an explanation. Namely, the branching ratio for the decay mode $B^+ \to K^+ \eta'$ is measured to be almost six times bigger than the one for $B^+ \to K^+ \pi^0$, although the same basic $b \to s$ penguin mechanism is expected to drive both processes. Apparently, this mechanism may have a different appearance when instead of the flavour octet pion there is an (almost) flavour singlet $\eta'$ particle involved.

It is well known that some extraordinary properties of the $\eta'$ particle are related to the QCD anomaly. Therefore, a suggestion to explain the enhancement of the $B \to K \eta'$ amplitude by the QCD anomaly at first sight looks very intriguing [3].

Our preceding investigation in this direction found that the $b \to s \eta'$ amplitude in the hard gluon regime represents a well defined short distance (SD) mechanism [4], but of minor numerical importance for the process at hand. The inability of this SD mechanism to account for the measured amplitude invites us to explore here the complementary long-distance (LD) mechanism. There are conclusions in the literature [5] that the “singlet penguin” amplitude (in the language of SU(3) diagrammatic approach [6]) contributes substantially to the $B \to K \eta'$ enhancement. One of the purposes of this note is to investigate this from another angle than done previously, i.e. to identify the singlet penguin contribution and estimate it within the well-defined microscopic dynamical framework.

In order to deal with the low energy properties of some of the involved gluons and to calculate the physical $B \to K \eta'$ amplitude, we have to introduce an appropriate low-energy description. We will employ ideas based on the chiral quark model ($\chi$QM)[7, 8] which has been used to describe $K \to 2\pi$ decays [9, 10]. Here we will rely on the extension of such models, namely the heavy-light chiral quark model (HL$\chi$QM) [11] which has, as the $\chi$QM, turned out to be a convenient calculational tool for addressing soft-gluon contributions, that can be expressed in terms of gluon condensate effects [8, 9, 10, 11, 12]. One should note that when the quarks and soft gluons in the HL$\chi$QM are integrated out, one obtains standard heavy light chiral perturbation theory (HL$\chi$PT) [13]. The HL$\chi$QM has been applied to $B - \bar{B}$ mixing [14] and to decays of the type $B \to D \bar{D}$ [15]. It has also been applied to the decay mode $B \to D \eta'$ [16] which has similar aspects as the mode we consider in this paper. In principle, the HL$\chi$QM naturally accounts only for soft kaons in the final state. However, as we show in the next section,
assuming the general form of the relevant form factor enables one to perform extrapolation from the soft to the hard kaon case.

2 Penguin contribution to $B \to K \eta'$ in the heavy-light chiral quark model

The LD mechanism that we propose for the $B \to K \eta'$ decay is shown in Figure 1. It accounts for the contribution obtained when a soft gluon ($G$) is emitted from the $B \to K$ transition together with the virtual gluon ($G^*$) associated with the penguin $b \to s$ transition. In addition, one of the gluons ($G'$) in the $\eta'GG$ vertex in Figure 1 is also assumed to be soft, so that two soft gluons form a vacuum condensate. The remaining off-shell gluon from the $b \to sG^*$ penguin is propagating into the $\eta'$. Now, the $B \to KGG^*$ vertex denoted by a large circle in Figure 1 can be calculated in the soft $K$ limit within the HLχQM, as displayed in Figure 2.

To begin with it, let us recall that the involved vector current form factors in heavy quark physics are defined by

$$
\langle K|\bar{s}\gamma^\mu b|B\rangle = f_+(q^2)(p_B + p_K)^\mu + f_-(q^2)(p_B - p_K)^\mu \\
= F_B(q^2)M_Bv^\mu + F_K(q^2)p_K^\mu ,
$$

(1)

where

$$
F_{B,K} \equiv f_+ \pm f_- , \quad \text{and} \quad p_B^\mu = M_Bv^\mu .
$$

(2)

In the $M_B \to \infty$ limit, $F_K$ dominates, as it can be seen from the scaling properties $F_K \sim \sqrt{M_B}$ and $F_B \sim 1/\sqrt{M_B}$ [17]. In addition, considering the soft $K$ limit within heavy-light chiral perturbation theory (HLχPT) [13], $F_K$
is dominated by the $B_s^*$ pole, and is given by

$$ (F_K)_{\text{soft}} = C_\gamma \frac{\sqrt{M_B} g_A \alpha_H}{f_\pi \sqrt{2} v \cdot p_K}, $$

(3)

where $\alpha_H = f_B \sqrt{M_B} / (C_\gamma + C_v)$, and $C_\gamma \approx 1$, $C_v \approx 0$ are the coefficients determined by QCD renormalization of the weak heavy-light current [18].

The soft $K$ limit is of course unphysical in our case, and to overcome this we employ a double pole structure

$$ F_K(q^2) = F_K(0) \frac{1}{(1 - \frac{q^2}{M_1^2})(1 - \frac{q^2}{M_2^2})}, $$

(4)

Such a structure, proposed in [19], seems to fit very well the existing data and the theoretical requirements on the heavy-light vector form factor, for, say, $M_1 = M_{B_s^*}$, and some parameter $\gamma$ fitting $M_2^2 = \gamma M_{B_s^*}^2$. To determine $\gamma$ we observe that (4) in our limits reads

$$ F_K(q^2)_{\text{soft}} = \frac{\gamma M_B F_K(0)}{2(\gamma - 1)} \frac{1}{v \cdot p_K}, $$

(5)

so that a comparison with (3) gives

$$ \frac{\gamma F_K(0)}{2(\gamma - 1)} = \frac{C_\gamma}{C_\gamma + C_v} \frac{f_B}{f_\pi \sqrt{2}} g_A, $$

(6)

which is general within HL$\chi$PT. Knowing the value for $F_K(0)$, $\gamma$ can be determined. We will use the result of the QCD sum rules on the light-cone analysis, $F_K(0) = 0.34 \pm 0.05$ [20], implying $\gamma = 1.27 \pm 0.08$, in agreement with lattice fits [19]. Thus, extrapolating from the soft $K$ to the general case we obtain the substitution rule

$$ \frac{1}{v \cdot p_K} \rightarrow \frac{2(\gamma - 1)}{M_B \gamma} \frac{1}{(1 - \frac{q^2}{M_1^2})(1 - \frac{q^2}{M_{B_s^*}^2})} = \frac{\sqrt{2} f_\pi F_K(q^2)}{f_B g_A \sqrt{M_B}}, $$

(7)

Below we will assume that the form factor for the $B \rightarrow KGG^*$ vertex also has the dipole form (4), because of $B_s^*$ pole dominance in both cases. This assumption will hold within HL$\chi$PT in the region where it is valid. Therefore we will adopt the rule (7) also for the $B \rightarrow KGG^*$ form factor. For this case, the position of the second pole might be somewhat different, which means that $\gamma$ and $F_K(0)$ in (6) should be replaced by $\gamma_G$ and $F_G^K(0)$, respectively.

The $b \rightarrow sG^*$ penguin operator at the quark level is

$$ g_s G_P \bar{s} \gamma^\mu \gamma^a \bar{L} t^a (DG)^\mu_a, $$
where $D$ means a covariant derivative, and

$$G_P = \frac{G_F}{\sqrt{2} \frac{\pi}{4}} V_{tb}^* V_{ts} (F_1^t - F_1^c).$$  \hfill (9)

For the quantity $(F_1^t - F_1^c)$ we take the one loop result \[21\], $0.26 - \left( \frac{2}{3} \ln \frac{m_c^2}{M_W^2} \right) \approx -5.2$. This result might be slightly changed by perturbative QCD effects like in \[22\] for $s \to dG^*$, but we do not enter such details here.

Note that the contribution of the dipole penguin operator $F_2 m_b \bar{s} \sigma^{\mu\nu} R t^a b G^{a\mu\nu}$ is suppressed by the small form factor $F_2 \approx 0.2$, and we neglect it.

The bosonization of the coloured quark current in (8) with emission of an additional soft gluon is known \[14\],

$$\bar{q}_L \gamma^\mu L t^a Q_v \longrightarrow \ G_H g_s G_{\alpha\beta}^{a} \text{ Tr} \left\{ \xi^\dagger \gamma^\mu L H_v \left[ A_1 \sigma^{\alpha\beta} + A_2 \sigma^{\alpha\beta} \gamma \cdot v \right] \right\},$$  \hfill (10)

where

$$H_v = \frac{1 + \gamma \cdot v}{2} \left( B^*_\mu \gamma^\mu - i B_5 \gamma_5 \right)$$  \hfill (11)

is the heavy meson “superfield” \[13\] and $\xi = \exp(i \Pi / f_\pi)$, with $\Pi$ being the standard Goldstone boson $3 \times 3$ matrix field. Furthermore, $G_H$ is the meson-quark coupling given by $G_H^2 = 2m_\rho / f_\pi^2$, where $m = 0.250 \pm 0.025$ GeV is the constituent light quark mass, $\rho \approx 1.1$ is a hadronic parameter \[14\] and $g_A = 0.59 \pm 0.08$ is the axial coupling of Goldstone bosons to heavy mesons. Moreover, within the HL$\chi$QM we find

$$A_1 = -\frac{1}{8} \left( \frac{1}{8\pi} - i I_2 \right), \quad A_2 = -\frac{1}{8} i I_2,$$  \hfill (12)
and $I_2$ is a logarithmically divergent loop integral which is expressed in terms of $f_\pi$ and the gluon condensate [8, 9, 10, 11], as follows

$$f_\pi^2 = -4im^2N_cI_2 + \frac{1}{24m^2}\langle \frac{\alpha_s}{\pi}G^2 \rangle.$$  

(13)

It should be noted that the structure in (10) should be rather general, while $G_H$ and the explicit expressions for $A_{1,2}$ are model dependent.

Taking the vector $(B^\ast)$ part of the $H_v$ and connecting with the $B^\ast$ propagator in Figure 2, we obtain in the soft $K$ limit the amplitude

$$M(B \rightarrow KGG^\ast)_{\text{soft}} = \frac{i g_{A}\sqrt{M_B}}{f_\pi \sqrt{2}} G_H G_P g_s G^\alpha_{\lambda\beta} (DG)^{\nu}_{\mu}(p_K)_\nu (v \cdot p_K) \times \left\{ -(A_1 + A_2)\epsilon^{\nu\alpha\beta} + 2A_2 v^\nu\epsilon^{\lambda\alpha\beta} v_\lambda \right\}. \quad (14)$$

In order to obtain the general amplitude for $B \rightarrow KGG^\ast$ from this equation, we perform the substitution (7) for $(v \cdot p_K)$ in the denominator above with $F_K(q^2)$ replaced by $F_K^G(q^2)$.

Concerning the $\eta'G^*G'$ interaction, it has the general form already used in ref. [16],

$$V(G^* \rightarrow G'\eta') = -\frac{1}{2}F_{\eta'gg}(q^2)\delta^{a'c}\epsilon^{\rho\lambda\kappa\sigma}\epsilon^{\ast a'c}G^c_{\lambda\kappa}\delta^{a\sigma}, \quad (15)$$

and several groups [23, 24] calculated the form-factor $F_{\eta'gg}(q^2)$ in the perturbative QCD approach. Since this approach becomes unreliable for gluon momenta of $\sim 1$ GeV, we adopt a formula from [24] which interpolates between the perturbative QCD region and the anomaly value for zero momentum. This formula gives $F_{\eta'gg}(m_{\eta'}^2) = 1.55 \pm 0.40$ GeV$^{-1}$, where an error of 25 % has been allowed. Accordingly, together with the value of the gluon condensate, this is the major source of uncertainty in our result below. Taking now the vacuum expectation value of the two soft gluons

$$g_s^2G^a_{\lambda\beta}G^c_{\lambda\kappa} \rightarrow \frac{4\pi^2\delta^{ac}}{12(N_c^2 - 1)}(g_{a\lambda}g_{b\kappa} - g_{a\kappa}g_{b\lambda})\langle \frac{\alpha_s}{\pi}G^2 \rangle,$$  

(16)

we obtain the final amplitude

$$M(B \rightarrow \eta')(G^2) = \frac{\pi^2G_PG_H F^G_K(q^2) F_{\eta'gg}(q^2) \langle \frac{\alpha_s}{\pi}G^2 \rangle}{3f_B \sqrt{M_B}} M^2_B \left\{ A_2 - 3(A_1 + A_2) \right\}, \quad (17)$$

where $q^2 = m_{\eta'}^2$.  

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Numerically, with $\langle \alpha_s \pi G^2 \rangle = (0.32 \pm 0.02 \text{ GeV})^4$, $f_B = 0.18 \pm 0.03 \text{ GeV}$, and $F^G_K(q^2) = F_K(q^2)$ given by the numbers below Eq. (6), we get

$$|M(B \to K\eta')| = (8 \pm 3) \times 10^{-9} \text{ GeV}.$$ \hspace{1cm} (18)

It should be noted, that even if the uncertainty of $\gamma$ is increased to 50 % when replaced by $\gamma_G$, it has no significant impact on the form factor $F^G_K(q^2)$ at the physical point $q^2 = m_{\eta'}^2$, and thereby not on the final result (18). This means that our assumption for the form factor $F^G_K(q^2)$ should be rather sound.

Our result (18) should be compared to the experimental amplitude $|M(B \to K\eta')_{\text{exp}}| = (88 \pm 2) \times 10^{-9} \text{ GeV}$. Thus, according to our analysis, only of order 10% of the $B \to K\eta'$ rate enhancement can be ascribed to this gluonic creation of $\eta'$. Some additional mechanisms are necessary, such as constructive interference of amplitudes for creating $\eta'$ in $d\bar{d}$ and $s\bar{s}$ state [25, 26]. However, the gluonic mechanism studied here, with its distinctive flavour-singlet nature, seems to go in the direction of the result of the SU(3) symmetry analysis in [5] that shows a substantial singlet penguin contribution to $B \to K\eta'$ amplitude.

3 Discussion

Let us comment here on how our result fits into the existing accounts of joining the two-gluon configurations arising from the $b \to sG$ and $b \to sGG$ transitions with those inherent to the $\eta'$ particle, in particular on those relying on properties of the $\eta'$ particle that are related to the QCD anomaly. The attempts to explain some puzzling hadronic weak decays by QCD anomalies are well known: the $\Delta I = 1/2$ enhancement in $K \to \pi\pi$ by the trace anomaly [27] and the enhancement of $B \to \eta'X_s$ decay rates by the axial anomaly [3, 28, 29].

We entered such study by employing the fact that anomaly permeates all distance scales, that enables one to study the role of two-gluon anomalous configurations from an extreme SD to a truly LD regime. Our recent study [4] shows that there is merely a remnant, the anomaly tail, in the extreme SD case. It should be noted that this contribution is obtained as a two quark operator for $b \to s\eta'$ and is very different from the LD contribution presented in detail in the previous section.

In addition, we have also identified some other contributions which we have found to be negligible. For example, Fritzsch [28] has suggested that an effective interaction of the form

$$H_{\text{eff}} = a \alpha_s G_F \bar{s}L b_R G_{\mu\nu} \hat{G}^{\mu\nu}$$ \hspace{1cm} (19)
might contribute significantly. We will describe elsewhere [30] different perturbative and nonperturbative contributions to such an effective interaction stemming from anomalous two-gluon configurations. Note that already a rotation of an appropriate term from Simma and Wyler’s paper [31] to Fritzsch’s form enables one to read off a perturbative contribution to coefficient $a$ above, $a_{SW}^{\text{pert}} \approx 4 \times 10^{-4}$ GeV$^{-1}$, which gives an amplitude 3–4 times smaller than (18).

After observing a systematic suppression of the anomalous two-gluon contributions through all the distance scales [30], we are focusing in the present paper to LD nonperturbative gluon configurations that may be phenomenologically more relevant. Thereby we are considering the low energy contribution where a gluon condensate accounts for the emission of soft gluons. A priori, our calculation would only be valid in the unphysical case where the outgoing kaon is soft. However, one can extrapolate to the physical point by introducing a dipole form factor for the $B \rightarrow K GG^*$ transition (as for the standard $B \rightarrow K$ transition current). As a result we find a more significant contribution from this mechanism, which can account for $\sim 10\%$ of the measured amplitude.

This result has to be compared with findings of [26] that flavour singlet contributions to $B \rightarrow K \eta'$ may be marginal. However, due to quite large uncertainties in both amplitudes these two results are actually not inconsistent. Note that the major portion of the singlet penguin amplitude in [26] comes from the operators corresponding to singlet quark configurations forming $\eta'$ particle. Recently another analysis within SCET appeared [32]. This analysis concludes that the “singlet penguin” contribution is essential to understand the process $B \rightarrow K \eta'$. Unfortunately, a direct comparison between our treatment of one soft ($q^2 \sim 0$) and one semi-hard ($q^2 \sim m_\eta^2$) gluon and the one by SCET is difficult to perform. Anyway, our contribution corresponding to the gluonic configurations forming the gluon condensate is a novel one, and may significantly increase the role of the singlet penguin mechanism in the direction of the result based on the SU(3) symmetry analysis [5].

A number of authors already used the surprising $B \rightarrow K \eta'$ enhancement to infer on the contributions beyond the Standard model. However, at this stage we need first to consider possible contributions from the specific mechanisms within the Standard model, like the one presented here. This mechanism seems to provide an additional contribution to the singlet penguin topology, the understanding of which may be of importance for explaining the data on CP asymmetries in penguin dominated modes [2].
Acknowledgment

K. K. and I. P. gratefully acknowledge the support of the Norwegian Research Council and the hospitality of the Department of Physics in Oslo, as well as support of Croatian Ministry of Science, Education and Sport under the contract No. 0119261. J. O. E. is supported in part by the Norwegian research council and by the European Union RTN network, Contract No. HPRN-CT-2002-00311 (EURIDICE). K. K. and J. O. E. thank Jure Zupan for fruitful discussions about the treatment of the considered process within SCET.

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