New Dipole Penguin Contribution to $K \rightarrow \pi\pi$ decays

Jan O. Eeg$^a$, Krešimir Kumerički$^b$, and Ivica Picek$^b$

$^a$Department of Physics, University of Oslo, P.O.B. 1048 Blindern, N-0316 Oslo, Norway

$^b$Department of Physics, Faculty of Science, University of Zagreb, P.O.B. 331, HR-10002 Zagreb, Croatia

October 17, 2008

Abstract

We point out that the standard chromomagnetic penguin dipole operator has a counterpart corresponding to off-shell momenta for external quarks. By employing the chiral quark model, we show that this new dipole penguin operator has the same bosonisation as the standard $Q_6$ operator. Accordingly, this new operator enlarges by $\sim 5\%$ the referent $Q_6$ contribution, which gives the dominant contribution to the $\varepsilon'/\varepsilon$ and also gives an important contribution to the $\Delta I = 1/2$ amplitude.

PACS: 13.20.Eb; 12.15.Ji; 12.39.-x; 12.39.Fe; 11.30.E

Keywords: K mesons, Electroweak transitions, Quark models, Chiral Lagrangians, CP violation

1j.o.eeg@fys.uio.no; corresponding author
2kkumer@phy.hr
3picek@phy.hr
1 Introduction

The physics of the $K \to 2\pi$ decay within the Standard Model has been a great challenge. First, the $\Delta I = 1/2$ rule is only vaguely understood, and similarly, the CP-violating quantity $\varepsilon'/\varepsilon$ has been very difficult to estimate because of the inherent hadronic uncertainties. In electroweak decays of K-mesons one constructs an effective Lagrangian at the quark level. Thereby one uses the equations of motion. In QED it is known that the main part of the Lamb shift disappears if the equation of motion is used for the electron self-energy, warning us that it is a bound state, nonperturbative effect. What about the Lamb shift like effects in the QCD context of hadronic decays?

An early study by two of us [1,2] was undertaken in order to account for the off-shell effects in $K$-meson decays. In particular, in ref. [1] we considered only the CP-conserving $K \to 2\pi$ amplitude, and the CP-violating off-shell part has been assigned to “the waiting list of pieces to be included in the reevaluation of $\varepsilon'/\varepsilon$”. At roughly the same time Bertolini et al. [3] included in such a reevaluation the chromomagnetic $Q_{11}$ penguin dipole contribution to $\varepsilon'/\varepsilon$, and in the present Letter we find as appropriate to consider its off-shell counterpart.

From the very beginning there has been a close relation between the gluonic penguin operators and the attempts to predict the direct CP-violation parameter $\varepsilon'/\varepsilon$ of the $K \to 2\pi$ amplitude. At first, the gluonic penguin had pointed to a possibility of a sizable (direct) CP-violating $K \to 2\pi$ amplitude. However, there had been a turning point with the large value of the $t$-quark mass, that led to a substantial cancellation between the dominant (gluonic $Q_6$, and electroweak $Q_8$) penguin contributions. That called for investigation of other possible contributions, including the off-shell contribution at hand.

For a brief history of progress on evaluating $\varepsilon'/\varepsilon$ and a more complete list of references on evaluation of $\varepsilon'/\varepsilon$ we refer to the review in ref. [4], and for later work to ref. [5]. These references show the way in which $\varepsilon'/\varepsilon$ is structured over the contributions from the operators belonging to, by now, standard operator basis [6]. The relatively large value for $\varepsilon'/\varepsilon$ obtained by the Trieste-Oslo collaboration [7] turned out to be a successful prediction for the outcome of the subsequent Fermilab [8] and CERN [9] measurements, that together with previous results from NA31 [10] and E731 [11] experiments reported the world average

$$\varepsilon'/\varepsilon = (16.7 \pm 1.6) \cdot 10^{-4}. \quad (1)$$

Basically, the prediction by the Trieste-Oslo group resides in the use of the chiral quark model ($\chi$QM) [12,13,14] which enabled to account for unavoidable nonperturbative QCD effects. Despite being a model approach, such
a treatment has some unique features: Besides weakening the destructive interference between the matrix elements of $Q_6$ and $Q_8$ operators (which we explicate in the next section), it enables the evaluation of the matrix elements for all the relevant operators within a single framework. Let us stress that in the chiral quark model we could evaluate the mentioned off-shell effects in K-meson decays \([1, 2, 15]\) and to account for the off-shell, off-diagonal self energy contribution to the CP-violating ratio $\varepsilon'/\varepsilon$ \([16]\). We believe that the off-shell chromomagnetic effect, calculated in this Letter, presents an interesting new piece in illuminating the $K \rightarrow \pi\pi$ puzzle.

2 Effective Lagrangian for $\Delta S = 1$ Decays

The amplitudes for $K \rightarrow 2\pi$ are described by an effective weak Lagrangian at quark level \([6, 17]\)

$$\mathcal{L}_W = \sum_i C_i(\mu) \, Q_i(\mu),$$

(2)

where all information on the short distance (SD) loop effects above a renormalisation scale $\mu$ is contained in Wilson coefficients $C_i$. These depend on the masses of the $W, Z$-bosons, the heavy quark masses ($m_q > \mu$), $\Lambda_{QCD}$ and on the renormalisation scheme. The $Q_i$’s are quark operators, typically containing products of two quark currents.

The standard basis (for $\mu < m_c$, relevant to kaon decays) includes ten 4-quark operators. We display the four most important for the $\Delta I = 1/2$ rule ($Q_{1,2,6}$) and $\varepsilon'/\varepsilon$ ($Q_{6,8}$):

$$Q_1 = 4 \left( \bar{s}L\gamma^\mu d_L \right) \left( \bar{u}L\gamma_\mu u_L \right), \quad Q_6 = -8 \sum_q (\bar{s}Lq_R) (\bar{q}Rd_L),$$

$$Q_2 = 4 \left( \bar{s}L\gamma^\mu u_L \right) \left( \bar{u}L\gamma_\mu d_L \right), \quad Q_8 = -12 \sum_q \hat{e}_q (\bar{s}Lq_R) (\bar{q}Rd_L),$$

(3)

where $\hat{e}_q$ are the quark charges ($\hat{e}_u = 2/3, \hat{e}_d = \hat{e}_s = -1/3$), and $q_{L,R}$ are the left- and right-handed projections of the quark fields.

Some studies also include the standard chromomagnetic dipole operator \([3, 4, 6, 18]\) that can be written as

$$Q_{11} = \frac{g_s}{8\pi^2} d \left[ m_s R + m_d L \right] \sigma \cdot G \, s + h.c.,$$

(4)

where $\sigma = \sigma^{\mu\nu} G_{\mu\nu}^a t^a$, $G_{\mu\nu}^a$ is the gluon field tensor, $t^a$ are the SU(3) generators normalised as $\text{Tr} (t^a t^b) = \delta^{ab}/2$, and $(L, R) = (1 \mp \gamma_5)/2$. However, we should stress that the operator for $s \rightarrow d + \text{gluon}(s)$ transition generated
by loop diagrams \[1, 19\] is not \(Q_{11}\) in eq. (4), but is given by

\[
\mathcal{L}(s \rightarrow dG) = B e^{\mu \nu \lambda \rho} G_{\mu \nu} (\bar{d}_L t^a i \bar{D}_\lambda \gamma_\rho s_L) \\
\rightarrow -\frac{1}{2} B \bar{d}_L [i \gamma \cdot D \sigma \cdot G + \sigma \cdot G i \gamma \cdot D] s_L .
\]

(5)

Here the second line is obtained by using simple algebra of Dirac matrices, and the coefficient \(B \sim g_s G_F \lambda_{K,M}\) depends on the loop integration. It is convenient to rewrite (5) as a sum of an off-shell term \(\mathcal{L}_G\) and the chromo-magnetic moment term \(\mathcal{L}_\sigma\) \[1\] :

\[
\mathcal{L}(s \rightarrow dG) = \mathcal{L}_G + \mathcal{L}_\sigma ; \\
\mathcal{L}_G = C_G Q_G ; \mathcal{L}_\sigma = C_{11} Q_{11} ,
\]

(6)

where we have introduced a counterpart of the standard dipole operator in eq. (4) :

\[
Q_G = \frac{g_s}{8 \pi^2} \bar{d} \left( (i \gamma \cdot D - m_d) \sigma \cdot G L + \sigma \cdot G R (i \gamma \cdot D - m_s) \right) s + h.c. .
\]

(7)

This operator vanishes by the QCD equation of motion for perturbatively interacting quark fields. The coefficients \(C_G\) and \(C_{11}\) above, being equal at the \(W\)-scale, evolve differently down to the scale \(\sim 1\) GeV, where hadronic matrix elements are evaluated. Therefore, in the next section we consider SD QCD corrections to the Wilson coefficient \(C_G\), that to our knowledge are not given in the literature.

In order to keep and calculate the contributions from \(Q_G\), one needs a framework to incorporate the effects of off-shell quarks at low energies, or equivalently, a framework where the operator \(Q_G\) cannot be rotated away. An important point of this Letter is that the operator \(Q_G\), to leading order, has the same bosonisation as the \(Q_6\) operator. This fact will enable us a direct comparison of the off-shell contribution coming from the operator \(Q_G\), and the leading CP-violating and CP-conserving contributions stemming from \(Q_6\).

The Lagrangian (6) is just a part of the more complete effective weak Lagrangian at quark level relevant to \(K\)-decays

\[
\mathcal{L}_W = \mathcal{L}_{4q} + \mathcal{L}_\sigma + \mathcal{L}_G + \mathcal{L}^{R}_{sd} ,
\]

(8)

where the additional terms are \(\mathcal{L}_{4q}\) corresponding to standard four quark operators \(3\) and \(\mathcal{L}^{R}_{sd}\), the renormalised off-diagonal self-energy \[1, 19\]

\[
\mathcal{L}^{R}_{sd} = -A \bar{d} (i \gamma \cdot D - m_d) i \gamma \cdot DR + M_R R + M_L L) (i \gamma \cdot D - m_s) s .
\]

(9)
The most important part of this last term for $\varepsilon'/\varepsilon$, the so-called self-penguin, was critically examined in ref. [20] and was shown to have a considerable off-shell contribution in $\chi$QM framework of ref. [16].

In the standard SD procedure [6, 17], $L_G$ and $L_{ds}^R$ in $L_W$, eq. (8), would be absent when applying the equation of motion at quark level. Instead, a more appropriate procedure should be to transform these terms away by a field redefinition, introducing new quark fields

$$
\begin{align*}
    d' &= d + B\sigma \cdot G L s - \frac{1}{2} A (i\gamma \cdot DR + M_R R + M_L L) (i\gamma \cdot D - m_s) s,
    \\
    s' &= s + B^* \sigma \cdot G L d - \frac{1}{2} A^* (i\gamma \cdot DR + M_R L + M_L R) (i\gamma \cdot D - m_s) d.
\end{align*}
$$

(10)

Then the parts $L_G$ and $L_{ds}^R$ involving the covariant derivatives are apparently removed, absorbed in the Dirac Lagrangian $L_f(q)$ for $q = (u, d, s)$ [1]:

$$
L_f(q) + L_G + L_{ds}^R = L_f(q'),
$$

(11)

where $L_f(q')$ is given later in eq. (24), with $q$ replaced by $q'$. In a strict SD treatment, primed and unprimed quark fields are equivalent. This means that $L_G$ (and similarly $L_{ds}^R$) does not contribute to $s \rightarrow d$ transitions for on-shell external quarks. In section 4 we will show how the effects of the mentioned off-shell operators reappear when low-energy strong interactions are taken into account in terms of the $\chi$QM.

3 \hspace{1em} The Wilson coefficient $C_G$

It is convenient to distinguish the CP-conserving and CP-violating parts of the Wilson coefficients for the $\Delta S = 1$ quark operators in (2). At some scale $\mu$ they can be written as

$$
C_i(\mu) = -\frac{G_F}{\sqrt{2}} \left[ \lambda_u z_i(\mu) - \lambda_t y_i(\mu) \right],
$$

(12)

where $G_F$ is the Fermi coupling, the functions $z_i(\mu)$ and $y_i(\mu)$ are the CP-conserving and CP-violating parts of the coefficients, respectively, and $\lambda_q = V_{qd} V_{qs}^*$ for $q = u, t$ are the CKM factors. The numerical values of $z_i$ and $y_i$ are in the range of order one down to $10^{-4}$, and can be found in the literature [6] for operators up to $Q_{11}$. In what follows we calculate the corresponding values for the operator $Q_G$, restricted to the truncated basis given by $Q_{\pm} = (Q_2 \pm Q_1)/2$ and $Q_G$ operators. Thereby we denote by $a_{\pm}$ the
Wilson coefficients of the 4-quark operators $Q_{\pm}$ with diagonal anomalous dimension matrix. In the logarithmic approximation, for $\mu \leq m_c$ they are

$$a_{\pm}(\mu^2) = \left[ \frac{\alpha_s(m_c^2)}{\alpha_s(\mu^2)} \right]^{\frac{d_+}{b(N_f)}} \left[ \frac{\alpha_s(\mu_c^2)}{\alpha_s(\mu_c^2)} \right] \left[ \frac{\alpha_s(M_W^2)}{\alpha_s(\mu_c^2)} \right]^{\frac{d_+}{b(N_f)}} \alpha_s(\mu_c^2) \alpha_s(\mu_c^2),$$  \hspace{1cm} (13)

where $d_+ = +2$ and $d_- = -4$ are the anomalous dimensions and $b(N_f) = 11N_c/3 - 2N_f/3$, where $N_f$ is the number of active flavours. The coefficient $a_{\pm}$ in the equation above, will become either the function $z(\mu)$ for the CP conserving part, or the function $y(\mu)$ for the CP violating part. For these cases, the values of $\mu$ will be $\mu \simeq 1$ GeV or $\mu = m_c$, respectively.

The anomalous dimension matrix for the truncated basis of the three operators $Q_{\pm}$ and $Q_n$ has, for $n = 6, 11, G$ the form:

$$\gamma = \begin{bmatrix} d_+ & 0 & X_+ \\ 0 & d_- & X_- \\ 0 & 0 & Y_n \end{bmatrix},$$  \hspace{1cm} (14)

where

$$X_+ = \begin{cases} \frac{11N_c}{18} - \frac{29}{18N_c} \pm \frac{3}{2} & \text{for } n = 11, G \\ \frac{1}{6} & \text{for } n = 6 \end{cases}$$

$$Y_n = \begin{cases} -6 \left( \frac{N^2-1}{2N_c} \right) + \frac{N_f}{3} & \text{for } n=6 \\ -2N_c - \frac{4}{N_c} & \text{for } n=11 \\ 0 & \text{for } n=G \end{cases}$$  \hspace{1cm} (15, 16)

One should note that $Q_G$, being an operator that vanishes by the perturbative QCD equation of motion, has zero anomalous dimension, $Y_G = 0$, in contrast to $Q_6$ and $Q_{11}$.

For handling the leading QCD corrections, there is a suitable prescription introduced in Refs. [21, 22] and applied by others [23, 24, 25, 26]. Using this prescription, one can write the amplitude as an integral over virtual quark loop momenta. The QCD-corrected coefficients $z_G$ and $y_G$ can be expressed in the integral form, which for CP-conserving case reads

$$z_G = \int_{\mu_c^2}^{M_W^2} \frac{dp^2}{p^2} \left( \frac{\alpha_s(p^2)}{4\pi} \right) \frac{1}{2} \left[ X_+ a_+(p^2) + X_- a_-(p^2) \right].$$  \hspace{1cm} (17)

For the CP-violating case we have

$$y_G = (F_2^c - F_2^t) + \int_{m_c^2}^{M_W^2} \frac{dp^2}{p^2} \left( \frac{\alpha_s(p^2)}{4\pi} \right) \frac{1}{2} \left[ X_+ a_+(p^2) + X_- a_-(p^2) \right],$$  \hspace{1cm} (18)
Table 1: Wilson coefficients of relevant operators obtained by leading order renormalisation group evolution from scale $\mu = M_W$, with flavour number thresholds at $m_b = 4.4$ GeV and $m_c = 1.3$ GeV, using $\alpha_s(M_Z) = 0.117$.

| Operator | $Q_6$ | $Q_{11}$ | $Q_G$ |
|----------|-------|----------|-------|
| $z(\mu = 1 \text{ GeV})$ | -0.009 | -0.033 | -0.035 |
| $y(\mu = m_c)$ | -0.083 | -0.318 | -0.407 |

where $F_2 \equiv F_2(m_t^2/M_W^2)$ are well-known Inami-Lim functions \cite{27}. Similarly, repeating the standard renormalisation group procedure of ref. \cite{6,28}, gives for the Wilson coefficient of $Q_G$, for the CP-conserving case,

$$ z_G(\mu) = \frac{X_+}{d_+} \left( \eta \hat{a}_+ - 1 \right) a_+(m_c) + \frac{X_-}{d_-} \left( \eta \hat{a}_- - 1 \right) a_-(m_c) , \hspace{1cm} (19) $$

where $\eta = \alpha_s(m_c)/\alpha_s(\mu)$, and $\hat{d}_\pm \equiv d_\pm/b(3)$. For the coefficient relevant for CP violation one obtains

$$ y_G(\mu = m_c) = (F_2^e - F_2^t) + \frac{X_+}{d_+} (a_+(m_c) - 1) + \frac{X_-}{d_-} (a_-(m_c) - 1) . \hspace{1cm} (20) $$

These expressions lead us to values for Wilson coefficient of $Q_G$ displayed in Table 1. In this Table we also give our values for coefficients of $Q_6$ and $Q_{11}$ that conform to those given in \cite{6}.

## 4 Bosonisation in the Chiral Quark Model

For light pseudoscalar mesons there is a well defined effective theory, chiral perturbation theory ($\chi PT$), having the basic symmetries of QCD. One can try to match $\chi PT$ to the weak Lagrangian at quark level, eq. \cite{2}, by *bosonizing* the quark operators $Q_i$:

$$ Q_i \rightarrow \sum_j F_{ij} \hat{L}_j , \hspace{1cm} (21) $$

where the $\hat{L}_j$'s are chiral Lagrangian terms having the symmetry of $Q_i$, and $F_{ij}$ are quantities to be calculated with non-perturbative methods (including quark models). Knowing the bosonization in \cite{21}, we could calculate the various $K$-decay amplitudes from a $\Delta S = 1$ chiral Lagrangian

$$ \mathcal{L}_W(\chi PT) = \sum_j G_j \hat{L}_j ; \hspace{1cm} G_j = \sum_i C_i F_{ij} . \hspace{1cm} (22) $$
The idea of such an approach is that the coefficients should be calculated (and matched) at the border of the SD and LD regimes. Thereby the factorized form of the coefficients $G_j$ in eq. (22) explicates the separation of SD contributions sitting in the $C_i$’s, and LD contributions residing in the $F_{ij}$’s.

In order to bosonise our relevant operators we employ the chiral quark model ($\chi$QM), that has been advocated by many authors [12, 13, 14] as an effective low-energy QCD. In this model, chiral-symmetry breaking is taken into account by adding a term to ordinary QCD:

$$L_{QCD} \rightarrow L_{QCD\chi} = L_{QCD} + L_\chi,$$

(23)

where $L_{QCD}$ in addition to the pure gluonic part contains the fermionic part

$$L_f(q) = \bar{q}(i\gamma \cdot D - M_q)q; \quad q = \begin{pmatrix} u \\ d \\ s \end{pmatrix},$$

(24)

Here $M_q = \text{diag} \left( m_u, m_d, m_s \right)$ is the current quark mass-matrix, whereas a non-perturbative term in (23), $L_\chi$, contains the parameter $m$ that is interpreted as the constituent quark mass ($\sim 200 - 250$ MeV). Note that the constituent and current masses are tied to different terms in the Lagrangian, with different transformation properties. Here $q$ is the SU(3) flavour triplet quark field, and $\Sigma$ contains the Goldstone-octet fields $\pi^a$:

$$\Sigma = \exp \left( i \sum_a \lambda^a \pi^a / f_\pi \right),$$

(26)

where $\lambda^a$ are the Gell-Mann matrices. The term $L_\chi$ contains meson-quark couplings. This means that the quarks can be integrated out and the coefficients of the various terms in the chiral Lagrangian are calculable from $L_{QCD\chi}$ in eq. (23).

Our procedure based on the equations (23, 24, 25) has to be understood in the following way: At scales above the cut-off $\Lambda_\chi$, the total Lagrangian is the sum of the standard $L_{QCD}$ and the weak effective Lagrangian $L_W$ in eq. (8), given in terms of the physical $u, d, s$-fields. Then at scales below $\Lambda_\chi$, the term $L_\chi$ in eq. (25) is turned on, so that the matrix elements of $L_W$ between mesonic states can be calculated owing to the meson-quark couplings.

The model has a “rotated” picture, where the term $L_\chi$ in (25) is transformed into a pure mass term $-m_\chi \chi$ for flavour rotated “constituent quark” fields $\chi_{L,R}$:

$$q_L \rightarrow \chi_L = \xi q_L \quad \text{and} \quad q_R \rightarrow \chi_R = \xi^\dagger q_R,$$

(27)
where $\xi \cdot \xi = \Sigma$. The meson–quark couplings in this rotated picture arise from the kinetic (Dirac) part of the constituent quark Lagrangian. These interactions can be described in terms of vector and axial vector fields coupled to constituent quark fields $\chi = \chi_R + \chi_L$. The sum of $\mathcal{L}_T$ in (24) and $\mathcal{L}_\chi$ in (25) are transformed into the equivalent form

$$\mathcal{L}_{\chi QM} = \bar{\chi} \left[ \gamma^\mu (iD_\mu + \mathcal{V}_\mu + \gamma_5 \mathcal{A}_\mu) - m \right] \chi - \bar{\chi} \tilde{M}_q \chi ,$$

where

$$\mathcal{V}_\mu = \frac{1}{2} \left[ \xi^\dagger (i\partial_\mu \xi) + \xi (i\partial_\mu \xi^\dagger) \right], \quad \mathcal{A}_\mu = \frac{1}{2} \left[ \xi^\dagger (i\partial_\mu \xi) - \xi (i\partial_\mu \xi^\dagger) \right] ,$$

and the current quark masses are residing in

$$\tilde{M}_q = \tilde{M}_q^V + \tilde{M}_q^A \gamma_5 , \quad \text{with} \quad \tilde{M}_q^{V(A)} = \frac{1}{2} (\xi^\dagger M_q \xi^\dagger \pm \xi M_q^\dagger \xi) .$$

Now, having quarks that are exposed to strong interactions, and are described by the chiral quark model at hand, we have to use the inverse of the transformation of eq. (10) into eq. (25). In this way the apparently removed terms are effectively reappearing as a new term in $\mathcal{L}_\chi$:

$$\mathcal{L}_\chi(q) = \mathcal{L}_\chi(q') + \Delta \mathcal{L}_\chi(q') + O(G_F^2) ,$$

where the new term $\Delta \mathcal{L}_\chi(q') \sim G_F$ introduces new vertices which compensate for those in $\mathcal{L}_G$ and $\mathcal{L}_{\tilde{d}S}$. When the fields $q'$ are integrated out, the result for a physical amplitude at mesonic level will be the same as if $\mathcal{L}_G$ and $\mathcal{L}_{\tilde{d}S}$ were applied without the field redefinition.

In general, each term in $\mathcal{L}_W$ which contains at least one power of $(i\gamma \cdot D - m_q)$ may be removed by a transformation like (10). However, for each term which is removed from $\mathcal{L}_W$, there will be a corresponding term appearing in $\mathcal{L}_\chi$. For the dipole operator $Q_G$, we obtain a contribution proportional to the constituent mass $m$

$$\Delta \mathcal{L}_\chi(q')_G = \frac{C_G}{8\pi^2} m \left( \overline{q'}_L \lambda_\sigma \cdot G \Sigma q'_R + \overline{q'}_R \Sigma^\dagger \lambda_\sigma \cdot G \lambda^\dagger \chi ' \right) ,$$

where $\lambda_\sigma = (\lambda_6 - i \lambda_7)/2$ is the combination of Gell-Mann matrices which transforms an $s$-quark into a $d$-quark. Employing the flavour rotation from eq. (27) we obtain the simple expression

$$\Delta \mathcal{L}_\chi(q')_G = \frac{C_G}{8\pi^2} \overline{\chi} F(-) \sigma \cdot G \chi' + h.c .$$
where
\[ F(-) = \xi \lambda \xi^\dagger. \] 
(34)
The expression (33) is ideal for bosonisation in terms of quark loops, whereas the analogous term for the self-energy still contains two derivatives, and can be calculated in a different way, as done in ref [16].

Using (28), the strong chiral Lagrangian \( O(p^2) \) can be understood as two axial currents \( A_\mu \) attached to a quark loop, leading to
\[ \mathcal{L}^{(2)}_s \sim \text{Tr} \left[ A_\mu A^\mu \right]. \] 
(35)
Using the relations
\[ 2i A_\mu = -\xi^\dagger (D^\mu \Sigma) \xi^\dagger = \xi (D^\mu \Sigma^\dagger) \xi, \] 
(36)
one obtains the leading strong chiral Lagrangian
\[ \mathcal{L}^{(2)}_s = \frac{f^2}{4} \text{Tr} \left( D^\mu \Sigma^\dagger D_\mu \Sigma \right), \] 
(37)
where \( D_\mu \) is the covariant derivative. Note that \( A_\mu \) is invariant under local chiral transformations [12, 13], in agreement with the invariance of \( \mathcal{L}^{(2)}_s \). In contrast, the vector field \( V_\mu \) transforms as a gauge field. Attaching in addition to two \( A_\mu \)'s also the mass term structures in (30), we will obtain the well-known \( L_5 \) term which enters the matrix element of \( Q_6 \).

In addition to the \( Q_6 \) operator, the referent object to which we compare our new off-shell dipole penguin is the chromomagnetic dipole operator eq. (39). It can be written in a chiral \( SU(3) \) invariant form, as a first step in its bosonisation procedure:
\[ Q_{11} = \frac{g_s}{8\pi^2} \left[ q_R M_q \lambda_- \sigma \cdot G q_L + \bar{q}_L \sigma \cdot G \lambda_- M^\dagger_q q_R \right], \] 
(38)
where \( q = (u, d, s) \). Note that this operator transforms as \( (8_L, 1_R) \) under the chiral \( SU(3)_L \times SU(3)_R \) symmetry if the current quark matrix is taken to transform as \( M_q \rightarrow V_R M_q V_L^\dagger \), where \( V_R \) and \( V_L \) are the chiral \( SU(3) \) transformation matrices.

In the next step, we write \( Q_{11} \) in the flavour rotated picture:
\[ Q_{11} = \frac{g_s}{8\pi^2} \left[ F^V_{(-)} + F^A_{(-)} \gamma_5 \right] \sigma \cdot G \chi, \] 
(39)
where \( F^V_{(-)} = \left( F^R_{(-)} \pm F^L_{(-)} \right) / 2 \) are expressed in terms of
\[ F^L_{(-)} = \xi^\dagger M_q \lambda_- \xi^\dagger, \quad \text{and} \quad F^R_{(-)} = \xi \lambda_- M^\dagger_q \xi. \] 
(40)
This operator is understood in terms of a quark loop. Let us first stress that its lowest-order contribution is \( \text{Tr} \left( F_{(-)} V \right) \), which is cancelled according to the FKW theorem \(^{29}\). To the NLO, we obtain a term corresponding to an interaction of \( F_{(-)} \) and two axial fields attached to a quark loop, shown in Fig. 1:

\[
\mathcal{L}^{(4)}(Q_{11}) \sim \text{Tr} \left[ F_{(-)} A_\mu A^\mu \right]. \quad (41)
\]

Note that \( F_{(-)} A \) is not contributing, because there must be an even number of \( \gamma_5 \)'s in the quark loop. Using (36) and (40), we find that \( \mathcal{L}^{(4)}(Q_{11}) \) can be written in the final bosonised form

\[
\mathcal{L}^{(4)}(Q_{11}) = G_8^{(4)}(Q_{11}) \text{Tr} \left[ (\Sigma^\dagger M_q \lambda_- + \lambda_- M_q^\dagger \Sigma) D^\mu \Sigma D_\mu \Sigma \right]. \quad (42)
\]

The coefficient \( G_8^{(4)}(Q_{11}) \) was calculated in \(^{18}\) to be \( \sim \langle \alpha_s \pi G^2 \rangle C_{11}/(8\pi^2) \), where the two-gluon condensate is a model dependent quantity in our approach. Finally, one can deduce \(^{18}\) the \( K \to 2\pi \) amplitude from the chiral structure of \( \mathcal{L}^{(4)}(Q_{11}) \):

\[
\mathcal{A}(K^0 \to \pi^+ \pi^-; Q_{11}) = \frac{\sqrt{2}}{f_\pi} (m_s - m_d) m_\pi^2 G_8^{(4)}(Q_{11}). \quad (43)
\]

Even though the coefficient in \( \mathcal{L}^{(4)}(Q_{11}) \) is large, the \( Q_{11} \) contribution to \( K \to 2\pi \) is small because of the factor \( m_\pi^2 \) (in place of \( m_K^2 \) obtained for \( Q_6 \)). Therefore, the modest role played by \( Q_{11} \) in \( \varepsilon'/\varepsilon \) is due to this kinematical suppression rather than due to being NLO in the chiral expansion.

The bosonisation of \( Q_6 \) follows basically the same line as for \( Q_{11} \) above. The standard expression for \( Q_6 \) obtained by Fierz transformation and displayed in (3) can be rewritten in the rotated picture as

\[
Q_6 = -8 (F_{(-)})_{\alpha\beta} (\chi_L)_\alpha (\chi_R)_\delta (\chi_R)_\delta (\chi_L)_\beta . \quad (44)
\]
where \( F(-) = \xi \lambda \xi^\dagger \), and Greek letters represent the flavour indices. Thus, likewise to \( \mathcal{L}^{(4)}(Q_{11}) \), the chiral representation of \( Q_6 \) to leading order can be written as
\[
\mathcal{L}^{(2)}(Q_6) \sim \text{Tr} \left[ F(-) A_\mu A^\mu \right].
\] (45)

By means of eq. (36) this can be written in the same form as other familiar \( \Delta S = 1 \) octet operators \( O(p^2) \) [30]
\[
\mathcal{L}^{(2)}(Q_6) = G_8^{(2)}(Q_6) \text{Tr} \left( \lambda_\Delta D^\mu \Sigma^\dagger D_\mu \Sigma \right).
\] (46)

This term gives rise to the \( K \to 2\pi \) amplitude
\[
\mathcal{A}(K^0 \to \pi^+ \pi^-; Q_6) = \frac{\sqrt{2}}{f^3} \left[ m_K^2 - m_\pi^2 \right] G_8^{(2)}(Q_6).
\] (47)

The relevant coefficient in this expression has been calculated [4] to be
\[
G_8^{(2)}(Q_6) = -16 \frac{L_5}{f^2} \langle \bar{q} q \rangle^2,
\] (48)
where \( L_5 \simeq 1.4 \times 10^{-3} \) is the coefficient of the \( O(p^4) \) chiral strong Lagrangian mentioned at the beginning of this section [4, 7, 18].

Finally, as indicated in eq. (33), the bosonisation of the operator \( Q_G \) proceeds by inserting the expression (33) as the shaded squares in Fig. 1. The corresponding loop evaluation gives
\[
G_8(Q_G) = -\frac{C_G}{8\pi^2} \left( \frac{\alpha_s}{\pi} G^2 \right).
\] (49)

For the gluon condensate in (49) we take the value \( \left( \frac{\alpha_s}{\pi} G G \right)^{1/4} = 310 \text{ MeV} \), in agreement, within uncertainties, with the lattice results and values used in refs. [7, 4, 31].

Both expressions (48) and (49) are concrete examples of the general separation of SD and LD effects, contained in the factorized form of the \( G_j \)'s in eq. (22) for our procedure of bosonisation. By performing the steps above, we have all ingredients that are necessary to estimate the new contribution of our dipole penguin operator \( Q_G \) on the same footing as the previously calculated contributions.

## 5 Results and Discussion

In the present Letter we are investigating a new dipole penguin operator that is the off-shell partner of the standard chromomagnetic penguin dipole
operator $Q_{11}$. Namely, in addition to the chromomagnetic penguin dipole operator \( (4) \) phrased \([6]\) as the mass insertions on external quark lines, there is, as explained in eqs. \([3]\) and \([4]\), an additional dipole operator $Q_G$ displayed in eq. \( (7) \) corresponding to off-shell momenta for external quarks confined in hadrons. Such Lamb-shift like effects in strong interactions are in another context very recently discussed in ref. \([32]\).

The new dipole penguin operator $Q_G$ studied here has several attractive features. We have shown that it has the same bosonisation as the standard $Q_6$ operator. Accordingly, $Q_G$ dominates over $Q_{11}$ which is higher order in chiral expansion. Indeed, the bosonised form $\mathcal{L}^{(4)}(Q_{11})$ results in the suppressed $K \rightarrow 2\pi$ amplitude $\mathcal{A}(K^0 \rightarrow \pi^+\pi^-; Q_{11})$ presented in eq. \( (13) \). Thereby, the bosonised form $\mathcal{L}^{(2)}(Q_6)$ leads to the referent $K \rightarrow 2\pi$ amplitude $\mathcal{A}(K^0 \rightarrow \pi^+\pi^-; Q_6)$ in eq. \( (47) \), and the two operators, $Q_G$ and $Q_6$, differ only in their respective coefficients, $G_8(Q_G)$ in eq. \( (49) \) and $G^{(2)}_8(Q_6)$ in eq. \( (48) \).

The ratio between $Q_G$ and $Q_6$ contributions can now be read from the LD hadronic factors and the SD Wilson coefficients contained in the coefficients $G_8(Q_G)$ and $G^{(2)}_8(Q_6)$:

$$\rho \equiv \frac{\mathcal{A}(K^0 \rightarrow \pi^+\pi^-; Q_G)}{\mathcal{A}(K^0 \rightarrow \pi^+\pi^-; Q_6)} = \frac{C_G}{C_6} h , \quad (50)$$

where the hadronic factor $h$ denotes the ratio of the respective LD pieces,

$$h = \frac{f_\pi^2 \langle \alpha_s \pi G^2 \rangle}{24 \cdot 8\pi^2 \cdot 16 L_5 \langle \bar{q}q \rangle^2} . \quad (51)$$

By substituting the numerical values, including $\langle \alpha_s \pi G G \rangle^{1/4} = 310$ MeV and $L_5 \simeq 1.4 \times 10^{-3}$, we obtain the hadronic factor $h \simeq 0.011$.

Finally, by employing the appropriate Wilson coefficients for the operators $Q_G$ and $Q_6$ given in Table 1, we obtain for the CP-violating and the CP-conserving parts of the ratio in eq. \( (50) \):

$$\rho_{\text{CP-violating}} = \frac{y_G}{y_6} h \simeq 0.05 , \quad (52)$$

$$\rho_{\text{CP-conserving}} = \frac{z_G}{z_6} h \simeq 0.04 . \quad (53)$$

This represents $\sim 5\%$ of the referent $Q_6$ contribution, to which it adds both in CP-violating and CP-conserving parts. In particular, there is a net coherent contribution from the CP-violating off-shell amplitudes, the one
from the new dipole operator considered here, and the previously calculated off-shell self energy contribution of \(\sim 15\%\) to \(\varepsilon'/\varepsilon\) in \([16]\). In conclusion, within the chiral quark model approach, we obtain in total an increase of \(\sim 20\%\) with respect to the leading \(Q_6\) contribution to the CP-violating ratio \(\varepsilon'/\varepsilon\) from off-shell operators. This result is still within the uncertainty of the theoretical value of Trieste-Oslo group, and although slightly higher than the world average \([1]\), it is closer to the new preliminary KTeV result \(\varepsilon'/\varepsilon = (19.2 \pm 1.1 \pm 1.8) \cdot 10^{-4}\) \([33]\).

Note added

After the first submission of this Letter we noticed another Letter \([32]\) addressing the effects of Lamb shift type in QCD. We also became aware of the new KTeV result based on doubling of the statistics and an improved control of systematics \([33]\).

Acknowledgement

This work was supported by the Ministry of Science, Education and Sport of the Republic of Croatia under contract No. 119-0982930-1016. J. O. E. thanks the Physics Department of University in Zagreb for its hospitality, and K. K. gratefully acknowledges the hospitality of the Department of Physics in Oslo. J. O. E. is supported in part by the Norwegian research council and by the European Commission RTN network, Contract No. MRTN-CT-2006-035482 (FLAVIAnet).

References

[1] J.O. Eeg and I. Picek, \textit{Off-shell s \(\rightarrow\) d transitions in K decays} Phys. Lett. \textbf{B301} (1993) 423.

[2] J. O. Eeg and I. Picek, \textit{New contribution to anomalous radiative K0 decays}, Phys. Lett. \textbf{B323} (1994) 193, \texttt{hep-ph/9312320}.

[3] S. Bertolini, M. Fabbrichesi and E. Gabrielli, \textit{The Relevance of the dipole penguin operators in \(\varepsilon'/\varepsilon\)} Phys. Lett. \textbf{B327} (1994) 136, \texttt{hep-ph/9312266}.
[4] S. Bertolini, M. Fabbrichesi and J. O. Eeg, *Theory of the CP-violating parameter $\varepsilon'/\varepsilon$*, Rev. Mod. Phys. 72 (2000) 65, [hep-ph/9802405], and references therein.

[5] J. Bijnens, Proceedings of Continuous Advances in QCD 2002, Minneapolis, 17-23 May 2002, p. 3 [hep-ph/0207082], and references therein; A. Pich, Proceedings of 32nd International Conference on High-Energy Physics (ICHEP 04), Beijing, China, 16-22 Aug 2004 [hep-ph/0410215], and references therein; J. Bijnens and J. Prades, *$\epsilon(K)/\epsilon(K)$ in the chiral limit*, JHEP 0006 (2000) 035, [hep-ph/0005189]; V. Cirigliano, J. F. Donoghue, E. Golowich, *Dimension-eight operators in the weak OPE*, JHEP 0010 (2000) 048, [hep-ph/0008281]; V. Cirigliano, G. Ecker, H. Neufeld, A. Pich, *Isospin breaking in $K \to \pi\pi$ decays*, Eur. Phys. J. C33 (2002) 369, [hep-ph/0310351].

[6] G. Buchalla, A. J. Buras and M. E. Lautenbacher, *Weak decays beyond leading logarithms*, Rev.Mod.Phys. 68 (1996) 1125, [hep-ph/9512380], and references therein.

[7] S. Bertolini, J. O. Eeg, M. Fabbrichesi and E. I. Lashin, *$\varepsilon'/\varepsilon$ at $O(p^4)$ in the chiral expansion*, Nucl. Phys. B514 (1998) 93, [hep-ph/9706260].

[8] A. Alavi-Harati et al. [KTeV Collaboration], *Observation of direct CP violation in $K_{S,L} \to \pi\pi$ decays*, Phys. Rev. Lett. 83 (1999) 22, [arXiv:hep-ex/9905060]; A. Alavi-Harati et al. [KTeV Collaboration], *Measurements of direct CP violation, CPT symmetry, and other parameters in the neutral kaon system*, Phys. Rev. D67 (2003) 012005 [Erratum-ibid. D70 (2004) 079904], [arXiv:hep-ex/0208007].

[9] V. Fanti et al. [NA48 Collaboration], *A new measurement of direct CP violation in two pion decays of the neutral kaon*, Phys. Lett. B465 (1999) 335, [arXiv:hep-ex/9909022]; A. Lai et al. [NA48 Collaboration], *A precise measurement of the direct CP violation parameter $Re(\varepsilon'/\varepsilon)$*, Eur. Phys. J. C22 (2001) 231 [arXiv:hep-ex/0110019]; J. R. Batley et al. [NA48 Collaboration], *A precision measurement of direct CP violation in the decay of neutral kaons into two pions*, Phys. Lett. B544 (2002) 97 [arXiv:hep-ex/0208009].

[10] G. D. Barr et al. [NA31 Collaboration], *A New measurement of direct CP violation in the neutral kaon system*, Phys. Lett. B317 (1993) 233.
[11] L. K. Gibbons et al., Measurement of the CP violation parameter Re $(\varepsilon'/\varepsilon)$, Phys. Rev. Lett. 70 (1993) 1203.

[12] S. Weinberg, Phenomenological Lagrangians, Physica 96A (1979) 327; A. Manohar and H. Georgi, Chiral Quarks and the Nonrelativistic Quark Model, Nucl. Phys. B234 (1984) 189; A. Manohar and G. Moore, Anomalous Inequivalence Of Phenomenological Theories, Nucl. Phys. B243 (1984) 55; S. Weinberg, The Axial vector coupling of the quark, Phys. Rev. Lett. 67 (1991) 3473.

[13] J. Bijnens, H. Sonoda and M.B. Wise, Power counting in $1/N_c$ in the chiral soliton model, Can.J.Phys. 64 (1986) 1; D.I. Diakonov, V.Yu. Petrov and P.V. Pobylitsa, A chiral theory of nucleons, Nucl. Phys. B306 (1988) 809; D. Espriu, E. de Rafael and J. Taron, The QCD effective action at long distances, Nucl. Phys. B345 (1990) 22; J. Bijnens, The Anomalous sector of the QCD effective Lagrangian, Nucl. Phys. B367 (1991) 709.

[14] A. Pich and E. de Rafael, Four quark operators and nonleptonic weak transitions, Nucl. Phys. B358 (1991) 311.

[15] D. Kekez, D. Klubuˇ car, K. Kumeriˇ cki and I. Picek, Quark off-shell contributions to $K_L \rightarrow \gamma\gamma$ in a bound state approach, Phys. Lett. B348 (1995) 239, [hep-ph/9501350].

[16] A. E. Bergan and J. O. Eeg, The self-penguin contribution to $K \rightarrow 2\pi$, Phys. Lett. B390 (1997) 420, [hep-ph/9609262].

[17] M. Ciuchini, E. Franco, G. Martinelli,and L. Reina, The $\Delta S = 1$ effective Hamiltonian including next-to-leading order QCD and QED corrections, Nucl.Phys. B415 (1994) 403, [hep-ph/9304257].

[18] S. Bertolini, M. Fabbrichesi and J. O. Eeg, Studying $\varepsilon'/\varepsilon$ in the chiral quark model: $\gamma_5$ scheme independence and NLO hadronic matrix elements, Nucl. Phys. B449 (1995) 197.

[19] H. Simma and D. Wyler, Hadronic rare B decays: the case $b \rightarrow sgg$, Nucl. Phys. B344 (1990) 283.

[20] B. Guberina, R. D. Peccei and I. Picek, A critique of Shabalin’s selfpenguin approach to the $\Delta I = 1/2$ rule, Phys. Lett. B188, (1987) 258.

[21] V. A. Novikov, M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, On the strong interaction effects on the $K_L \rightarrow 2\mu$ decay and $K_L-K_S$ mass difference. A Reply, Phys. Rev. D16 (1977) 223.
[22] A.I. Vainstein, V.I. Zakharov, L.B. Okun’ and M. A. Shifman, On the $K^+ \to \pi^+ e^+ e^-$ decay, Yad. Fiz. 24 (1976) 820, [Sov. J. Nucl. Phys. 24 (1976) 427].

[23] W.A. Bardeen, A.J. Buras and J.-M. Gerard, The $K \to \pi \pi$ decays in the large $n$ limit: quark evolution, Nucl. Phys. B293 (1987) 787; A. Pich and E. de Rafael, Four quark operators and nonleptonic weak transitions, Nucl. Phys. B358 (1991) 311.

[24] J.O. Eeg, Short distance QCD effects on the $s \to d \gamma^*$ loop, Phys. Rev. D23 (1981) 2596.

[25] J.O. Eeg, B. Nižić and I. Picek, Direct CP violation in $K_L, K_S \to 2 \gamma$ for large top mass, Phys. Lett. B244 (1990) 513.

[26] J.O. Eeg, K. Kumerički and I. Picek, On the short-distance dispersive two-photon $K_L \to \mu^+ \mu^-$ amplitude, Eur. Phys. J. C1 (1998) 531, [hep-ph/9605337].

[27] T. Inami and C. S. Lim, Effects of superheavy quarks and leptons in low-energy weak processes $K_L \to \mu \bar{\mu}, K^+ \to \pi^+ \nu \bar{\nu}$ and $K^0 \leftrightarrow \bar{K}^0$, Prog. Theor. Phys. 65 (1981) 297, [Erratum-ibid. 65 (1981) 1772].

[28] B. Grinstein, R. Springer and M. B. Wise, Effective Hamiltonian for weak radiative $B$ meson decay, Phys. Lett. B202 (1988) 138.

[29] G. Feinberg, P. Kabir and S. Weinberg, Transformation of muons into electrons, Phys. Rev. Lett. 3 (1959) 527; J.F. Donoghue and B.R. Holstein, Reanalysis of higgs boson exchange models of CP violation, Phys. Rev. D32 (1985) 1152.

[30] J.A. Cronin, Phenomenological model of strong and weak interactions in chiral $U(3) \times U(3)$, Phys. Rev. 161 (1967) 1483.

[31] A. Hiorth and J.O. Eeg, A heavy light chiral quark model, Phys. Rev. D66 (2002) 074001.

[32] S.J. Brodsky and R. Shrock, Maximum wavelength of confined quarks and gluons and properties of quantum chromodynamics, Phys. Lett. B666 (2008) 95, [arXiv:0806.1535 [hep-th]].

[33] Talk by M. Sozzi at Workshop on the origins of P, CP and T violation, ICTP, Trieste, 2-5 July 2008.