Charming penguin in nonleptonic $B$ decays

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ABSTRACT

In the study of two-body charmless $B$ decays as a mean of looking for direct CP-violation and measuring the CKM mixing parameters in the Standard Model, the short-distance penguin contribution with its absorptive part generated by charm quark loop seems capable of producing sufficient $B \to K\pi$ decays rates, as obtained in factorization and QCD-improved factorization models. However there are also long-distance charming penguin contributions which also give rise to a strong phase due to the rescattering $D^*D^* \to K\pi$ . In this talk, I would like to discuss \cite{19} a recent work on the long-distance charming penguin as a a different approach to the calculation of the penguin contributions in $B \to K\pi$ decays from charmed meson intermediate states. Using chiral effective Lagrangian for light and heavy mesons, corrected for hard pion and kaon momenta, we show that the charming-penguin contributions increase significantly the $B \to K\pi$ decays rates from its short-distance contributions, giving results in better agreement with experimental data.

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1 Introduction

Recent measurements by the CLEO \cite{1}, Babar \cite{2} and Belle \cite{3} collaboration give consistent values for the $B$-$K\pi$ branching ratios, which are respectively $(18.2^{+4.6}_{-4.0} \pm 1.6) \times 10^{-6}$, $(18.2^{+3.3}_{-3.0} \pm 2.0) \times 10^{-6}$, $(13.7^{+5.7}_{-4.8} \pm 1.9) \times 10^{-6}$ for $B^+\to K^0\pi^+$ and $(17.2^{+2.5}_{-2.4} \pm 1.2) \times 10^{-6}$, $(16.7 \pm 1.6^{+1.2}_{-1.7}) \times 10^{-6}$, $(19.3^{+3.4}_{-3.2} \pm 0.6) \times 10^{-6}$ for $B^0\to K^+\pi^-$ decays. The short-distance contributions to $B \to K\pi$ decays as given by the penguin operators without charm quark loop in factorization model seem to produce the $B \to K\pi$ decays rates too small compared to the data \cite{4}. A better agreement is obtained by including the so-called charming penguin contribution in the effective Wilson coefficients \cite{5, 6, 7, 8}. In this way an absorptive part of the decay amplitude is generated and the strong phase from this absorptive part can produce CP violation in $B \to K\pi$ decays \cite{5, 10}. This approach seems to produce decay rates in agreement with data, at least qualitatively, as shown previously \cite{4, 11, 12, 13}, where the charm quark loop contribution increases the effective Wilson coefficients of the strong penguin operators by about 30\%. More recently charm quark effects computed by this method have been obtained in recent works dealing with the validity of factorization \cite{14, 15, 16}. Another approach is to assume that the charm quark contributions are basically long-distance effects essentially due to rescattering processes such as, e.g. $B \to DD_s \to K\pi$. These contributions, first discussed in \cite{17}, have been more recently stressed by \cite{4}, where they are called charming penguin terms. The situation is similar to the $B_s \to \gamma\gamma$ decay for which the absorptive part obtained in \cite{18} is comparable to the short-distance contribution. I would like to discuss here a recent work \cite{19} on the charming penguin contributions in $B \to K\pi$ decays. As details can be found in this reference, I will present only the main results of the work.

2 Short and Long distance weak matrix element

In the standard model, effective Hamiltonian for non-leptonic $B$ decays are given by

$$
\mathcal{H}_{\text{eff}} = \frac{G_F}{\sqrt{2}} \left[ V_{ub}^* V_{us} (c_1 O_1^u + c_2 O_2^u) + V_{cb}^* V_{cs} (c_1 O_1^c + c_2 O_2^c) - V_{tb}^* V_{ts} \left( \sum_{i=3}^{10} c_i O_i + c_g O_g \right) \right] (1)
$$

where $c_i$ are the Wilson coefficients evaluated at the normalization scale $\mu = m_b$ \cite{3, 6, 21, 22} and next-to-leading QCD radiative corrections are included. $O_1$ and $O_2$ are the usual tree-level operators, $O_i$ ($i = 3, \ldots, 10$) are the penguin operators and $O_g$ is the chromomagnetic gluon operator.
The $B \to K\pi$ decay amplitude $A_{K\pi}$ is given by

$$A_{K\pi} = <K(p_K)\pi(p_\pi)|iH_{\text{eff}}|B(p_B)>.$$  \hspace{1cm} (2)

In the factorization approximation, the above matrix element is evaluated at the tree-level as higher order QCD radiative corrections are already included in the effective Wilson coefficients and the charm quark operators $O_1^c$ and $O_2^c$ do not contribute. The short-distance part $A_{SD}$ is obtained with $c_2 = 1.105$, $c_1 = -0.228$, $c_3 = 0.013$, $c_4 = -0.029$, $c_5 = 0.009$, $c_6 = -0.033$ \cite{20}; $|V_{ub}| = 0.0038$, $V_{us} = 0.22$, $V_{tb} \simeq 1$, $V_{ts} = -0.040$ and $\gamma = -\arg(V_{ub}) = 54.8^\circ$ \cite{23} and $F_{0}^{B \to \pi}(m_{K}^2) = 0.37$. We find

$$A_{SD}(B^+ \to K^0\pi^+) = 2.43 \times 10^{-8} \text{ GeV}$$
$$A_{SD}(B^0 \to K^+\pi^-) = (1.86 - i 0.95) \times 10^{-8} \text{ GeV}.$$ \hspace{1cm} (3)

As mentioned, the $B \to K\pi$ branching ratios obtained from Eq.(3) are too small compared with experiments. Instead of using perturbative QCD to treat the charm quark loop contributions, we now consider the one-particle $D, D^*$ intermediate state contribution to the T-product of two charged weak currents corresponding to the local operators $O_2^c$. The matrix element of $O_2^c$ is evaluated using a sum rule due to Wilson \cite{24}. Following Wilson, consider now the short-distance limit of the T-product of two weak currents

$$T[J_{\mu N}(x)J_{\mu S}(0)] = B_1'(x)\sigma'_m(0)$$ \hspace{1cm} (4)

where the contributions from the more singular, lower dimension operators have been taken out. $B_1'(x)$ is the coefficient of the local operator $\sigma'_m(0)$. Let $M_{AB}(q) = \int d^4x \exp(iq \cdot x) <A|J_{\mu S}(x)J_{\mu N}(0)|B>$ we have in momentum space,

$$\int^{q_{\text{max}}} M_{AB}(q)d^4q = B_1'(q_{\text{max}})\sigma'_{AB}, \quad B_1'(q_{\text{max}}) = \int^{q_{\text{max}}} B_1'(q)d^4q$$ \hspace{1cm} (5)

If $B_1'(x)$ scales as $(x^2)^0$ as in QCD, and for $q_{\text{max}}$ not too large, we obtain

$$\int^{q_{\text{max}}} M_{AB}(q)d^4q = \sigma'_{AB}$$ \hspace{1cm} (6)

Eq.(8) thus gives us the matrix element of the local operators in terms of a Cottingham-like formula evaluated only up to a cut-off momentum $q_{\text{max}}$ as the high momenta of the integral has already been factorized in the Wilson coefficients, as stressed in previous work \cite{17, 23, 26}. It should be stressed here that in factorization model, the exchange term in the effective Hamiltonian is usually Fierz-reordered into a product of two color-singlet operators and then evaluated by vacuum saturation. Actually, it can also be expressed in terms of an integral over the virtual momentum $q$ which is the difference of the two quark momenta in the initial and final hadron.
Table 1: Numerical values for the real part of $A_{LD}$ in GeV for $\mu_\ell = 0.5 - 0.7$ GeV. First column refers to the $D$, the second is the $D^*$ contribution.

| $\mu_\ell$ | $D$          | $D^*$         | Total         |
|------------|--------------|---------------|---------------|
| 0.5        | $-4.66 \times 10^{-9}$ | $1.62 \times 10^{-8}$ | $1.15 \times 10^{-8}$ |
| 0.6        | $-7.77 \times 10^{-9}$ | $2.79 \times 10^{-8}$ | $2.01 \times 10^{-8}$ |
| 0.7        | $-1.19 \times 10^{-8}$ | $4.40 \times 10^{-8}$ | $3.21 \times 10^{-8}$ |

For example, the exchange term in the $K\pi$ transition is given as $(\psi(k, k - p)$ is the pion B-S wave function),

$$A(K^- \to \pi^-) = \int \frac{d^4 k}{(2\pi)^4} \int \frac{d^4 k'}{(2\pi)^4} \bar{\psi}(k, k - p) T_W(k, k - p; k', k' - p') \psi(k', k' - p')$$  \(7\)

Making a change of variable $k' = q + k$, we have

$$A(K^- \to \pi^-) = \int \frac{d^4 q}{(2\pi)^4} T(p, q)$$  \(8\)

$$T(p, q) = \int \frac{d^4 k}{(2\pi)^4} \bar{\psi}(k, k - p) T_W(k, k - p; k + q, k + q - p) \psi(k + q, k + q - p)$$  \(9\)

which is a higher twist contribution to the forward virtual scattering of the $W$ boson with momentum $q$ off a hadron. A similar expression can also be given for the transition $\Sigma \to p$ in hyperon nonleptonic decays. The above expression shows that nonleptonic weak matrix elements can be expressed as integral over the virtual $W$ boson scattering amplitude. We have, for the long-distance part $A_{LD}$

$$A_{LD} = A_{LD}(B^+ \to K^0\pi^+) = A_{LD}(B^0 \to K^+\pi^-) = \frac{G_F}{\sqrt{2}} V_{cb}^* V_{cs} a_2 \int \frac{d^4 q}{(2\pi)^4} \theta(q^2 + \mu^2) T(q, p_B, p_K, p_\pi)$$  \(10\)

where $\mu$ (or $q_{max}$) is a cut-off momentum separating long-distance and short-distance contribution. $T(q, p_B, p_K, p_\pi) = g^{\mu\nu} T_{\mu\nu}$, with

$$T_{\mu\nu} = i \int d^4 x \exp(i q \cdot x) < K(p_K)\pi(p_\pi) | T(J_\mu(x)J_\nu(0)) | B(p_B) >$$  \(11\)

$J_\mu = \bar{b}\gamma_\mu(1 - \gamma_5)c$ and $J_\nu = \bar{c}\gamma_\nu(1 - \gamma_5)s$.

To compute $A_{LD}$ we saturate the $T_{\mu\nu}$ with the $D, D^*$ intermediate states. This gives us the usual $D, D^*$ pole term (Born term) for $T(q, p_B, p_K, p_\pi)$. To compute these pole terms, we use heavy quark effective theory and chiral effective lagrangian to obtain the $B \to D, D^*$ and $D \to K\pi$ and $D^* \to K\pi$ semi-leptonic decay form
factors \[27\] which appear at each vertex of the pole diagrams. < \(D, D^*\)|\(J^\mu|B\) is parameterized in terms of the Isgur-Wise function and the matrix elements < \(K\pi|J^\mu|D\) and < \(K\pi|J^\mu|D^*\) > are computed using Chiral Effective Lagrangian for semileptonic decays of heavy mesons to light pseudo-scalar mesons. We extrapolate the soft meson limit to higher momenta by using the full \(D^*\) propagator in the pole terms (a similar use of the full \(D^*\) propagator to go beyond the soft pion result has also been given in \[28\]) . We also introduce a form factor in the strong \(DD\pi\) coupling constant (a similar approach is used in semileptonic decays \[29\]). Including this effect, we obtain, for hard pion,

\[
G_{D^*D\pi} = \frac{2m_D}{f_\pi} \frac{g}{f_\pi} \frac{F(|\vec{p}_\pi|)}{\omega^2 - 1} ,
\]

where \(F(|\vec{p}_\pi|)\) is normalized by \(F(0) = 1\) which corresponds to the soft pion limit. (\(g \approx 0.4\)). This form factor can be evaluated by using the constituent quark model which gives roughly, for \(|\vec{p}_\pi| \approx m_B/2\), \(F(|\vec{p}_\pi|) = 0.065 \pm 0.035\).

Since the threshold for the \(D, D_s\) and \(D, D_s^*\) production is below the \(B\) meson mass, the \(D_s\) and \(D_s^*\) pole term for the \(D, D^* \rightarrow K\pi\) form factors have an absorptive part. This pole term is in fact a rescattering term via the Cabibbo-allowed \(B \rightarrow D, D_s^*\) decays followed by the strong annihilation process \(D, D_s^* \rightarrow K\pi\) and can be obtained from the unitarity of the \(B \rightarrow K\pi\) decay amplitude. We have

\[
\text{Disc } A_{LD} = 2i \text{ Im } A_{LD} = (-2\pi i)^2 \int \frac{d^4q}{(2\pi)^4} \delta^+(q^2 - m_{D_s}^2) \delta^+(p_{D_s^*}^2 - m_{D_s^*}^2) \times A(B \rightarrow D_s(s)D(s)) A(D_s(s)D(s) \rightarrow K\pi) =
\]

\[
= -\frac{m_D}{16\pi^2 m_B} \sqrt{\omega^2 - 1} \int d\tilde{n} A(B \rightarrow D_s(s)D(s)) A(D_s(s)D(s) \rightarrow K\pi) .
\]

With the \(A(B \rightarrow D_sD), A(B \rightarrow D_s^*D^*)\) given by factorization and \(A(D_sD \rightarrow K\pi), A(D_s^*D^* \rightarrow K\pi)\) by the \(t\)-channel \(D, D^*\) exchange pole terms which are proportional to \(G_{D^*D\pi}^2\) and could be large due to the factor \(m_D^2\). However the rescattering amplitudes \(A(D_s^*D^* \rightarrow K\pi)\) etc. which are exclusive processes at high energy, should be suppressed. This is taken account by the suppression factor \(F(|\vec{p}_\pi|)\) mentioned above.

We find, for the absorptive part

\[
\text{Im } A_{LD} = 2.34 \times 10^{-8} \text{ GeV} \quad (14)
\]

of which \(1.45 \times 10^{-8} \text{ GeV}\) and \(0.89 \times 10^{-8} \text{ GeV}\) are respectively the \(D, D_s\) and \(D^*, D_s^*\) contributions. To find the real part, we compute all Feynman diagrams obtained with the effective Lagrangian for the weak form factors and integrate over the virtual
current momentum $q$ up to a cut-off $\mu = m_b$. This includes the direct term and the pole terms which produce the absorptive part. It is possible to choose a cut-off momentum by a change of variable $q = p_B - p_{D^(*)}$ to the momentum $\ell$ defined by the formula

$$q = p_B - p_{D^(*)} \equiv (m_B - m_{D^(*)})v - \ell.$$  

As discussed in [19], the chiral symmetry breaking scale is about 1 GeV and the mean charm quark momentum $k$ for the on-shell $D$ meson is about 300 MeV, the virtual momentum $\ell$ should be below 0.6 GeV, hence a cut-off $\mu_{\ell} \approx 0.6$ GeV. The real part is then given by a Cottingham formula as follows [27]

$$\text{Re} \ A_{LD} = \frac{i}{2} \frac{G_F}{\sqrt{2}} V_{cb}^* V_{cs} a_2 \int_{0}^{\mu_{\ell}^2} dL^2 \int_{-\sqrt{L^2}}^{+\sqrt{L^2}} dl_0 \sqrt{L^2 - l_0^2} \int_{-1}^{1} d\cos(\theta) \ i \times \left\{ \frac{j_{D^*}}{p_{D^*}^2 - m_{D^*}^2} + \sum_{\text{pol}} \frac{j_{D_{s}^*}^*}{p_{D_{s}^*}^2 - m_{D_{s}^*}^2} \right\}. \quad (16)$$

in the above expression, the coupling constant $g$ are corrected by the form factor $F(|\vec{p}_\pi|)$. The results for the real part are shown in Table 1 for $\mu_{\ell} = 0.5 - 0.7$ GeV. Our numerical results show that the long-distance charming penguin contributions to the decays $B \to K\pi$ are significant. These results agree qualitatively with a phenomenological analysis of these contributions given in [4]. In particular, we found that the absorptive part due to the $D, D_{s}$ states is somewhat bigger than that from the $D^*, D_{s}^*$ states, but of the same sign. The real part due to the $D^*, D_{s}^*$ states is however 3 - 4 times bigger and opposite in sign to the contributions from the $D, D_{s}$ states. As shown in Table 1, the real part and absorptive part are of the same order of magnitude, at the $10^{-8}$ GeV level. The results for the branching ratios are

$$\mathcal{B}(B^+ \to K^0 \pi^+) = (2.4_{-1.9}^{+2.7}) \times 10^{-5}$$
$$\mathcal{B}(B^0 \to K^+ \pi^-) = (1.5_{-1.3}^{+1.8}) \times 10^{-5}. \quad (17)$$

which agrees with the results from CLEO [1], Babar [2] and Belle [3] mentioned above. The inelastic FSI strong phase we get from the absorptive part will produce a CP violation in $B \to K\pi$ decays via the interference with the tree-level terms. We get, for the CP-asymmetry between $B^0 \to K^+ \pi^-$ and $\bar{B}^0 \to K^- \pi^+$ decay rates : $A_{CP} = 0.21$ for $\gamma = 54.8^0$ which is comparable with recent results from CLEO [30].

### 3 Conclusion

In conclusion, we believe that the charmed resonance contributions we found seem to be capable of producing the charming-penguin terms suggested in [4] within theoretical errors. The strong phase generated by the real charm meson intermediate states
would be the essential mechanism for direct CP violation in charmless $B$ decays as suggested by [5, 10]. Though our estimate of the real part get uncertainties from the value of the cut-off momentum $\mu_\ell$ due to various form factors, its strength is comparable with the short-distance part, though not as important as the long-distance contribution in $K \to \pi\pi$ decays.

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