Importance of Dipole Penguin Operator in $B$ Decays *

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(September, 1995)

Abstract

The importance of dipole penguin operator, $O_{11} = \left(\frac{g_s}{8\pi^2}\right) m_b \bar{s} \sigma_{\mu\nu} R T^{ab} b G^\mu_{a\nu}$, on $B$ decays is demonstrated. It is shown that branching ratios for the inclusive decay $b \to s\bar{s}s$, the semi-inclusive decay $b \to s\phi$, and the exclusive decay $B \to \pi K$ are enhanced by 20% to 30%. A useful consequence of this calculation is the inclusive $b \to s\bar{s}s$ spectrum which will be essential in isolating pure penguin process.

*Work supported in part by the Department of Energy Grant No. DE-FG06-85ER40224.
Flavor changing one loop processes mediated by a gluon play an important role in understanding the Standard Model (SM) [1–7]. It is by now well known that substantial corrections arise in these estimates from inclusion of the full electroweak effects [2,4–7] (the $\gamma$, $Z$ penguins and the “box” contributions). Naively, one would think that in the $B$ system electroweak corrections to gluon exchange may be negligible. This is, in fact, incorrect. The large mass of the top quark increases the relative contributions of the electroweak corrections. Our recent estimates for $b \to s\phi$ process suggests a reduction in the rate of $20\% \sim 30\%$ [4,5]. In this paper we study another contribution, which has been mostly neglected so far, to nonleptonic $B$ decays. This is the contribution from the Dipole Penguin Operators (DPO). Similar operators have been shown to give non-negligible contribution to $\epsilon'/\epsilon$ [8,9] and hyperon decays [9]. We consider three different types of decays in increasing model dependence, the inclusive $b \to sq\bar{q}$ and $b \to ss\bar{s}$, the semi-inclusive $b \to s\phi$, and the exclusive $B^- \to K^0\pi^-$ decays. We find that the DPO gives important additional contributions to these decays and enhances their branching ratios by up to 30%. Any future quantitative analysis of $B$ decays that relies on full electroweak Hamiltonian will have to include the DPO contributions. An important consequence of this calculation is the inclusive $b \to ss\bar{s}$ spectrum which will be useful in isolating pure penguin process.

The effective Hamiltonian

The effective Hamiltonian up to one loop level in electroweak interaction for charmless $B$ decays can be written as

$$H_{\Delta B=1} = \frac{4G_F}{\sqrt{2}} [V_{ub}V_{us}^* (c_1 O_1 + c_2 O_2) - V_{tb}V_{tq}^* \sum_{i=3}^{12} c_i O_i] + H.C.,$$

where $O_i$ are defined as

$$O_1 = \bar{q}_\alpha \gamma_\mu L f_\beta \bar{f}_\gamma \gamma^\mu L b_\alpha, \quad O_2 = \bar{q}_\gamma L f \bar{f} \gamma^\mu L b,$$

$$O_{3(5)} = \bar{q}_\gamma L b \sum_{q'} \bar{q}_{\alpha}' \gamma_\mu L(R) q', \quad Q_{4(6)} = \bar{q}_\alpha \gamma_\mu L b_\beta \sum_{q'} \bar{q}_{\beta}' \gamma_\mu L(R) q'_\alpha,$$

$$O_{7(9)} = \frac{3}{2} \bar{q}_\gamma L b \sum_{q'} e_{q'} \bar{q}_{\alpha}' \gamma_\mu R(L) q', \quad Q_{8(10)} = \frac{3}{2} \bar{q}_\alpha \gamma_\mu L b_\beta \sum_{q'} e_{q'} \bar{q}_{\beta}' \gamma_\mu R(L) q'_\alpha,$$

$$O_{11} = \frac{g_s}{32\pi^2} m_b \bar{q}_\mu \sigma_{\mu\nu} R T_a b G^\mu_a, \quad Q_{12} = \frac{e}{32\pi^2} m_b \bar{q}_\mu \sigma_{\mu\nu} R b F^{\mu\nu},$$
where \( L(R) = (1 \mp \gamma_5)/2 \), \( f \) can be \( u \) or \( c \), \( q \) can be \( d \) or \( s \), \( q' \) is summed over \( u, d, s, \) and \( c \). \( \alpha \) and \( \beta \) are the color indices. \( T^a \) is the SU(3) generator with the normalization \( \text{Tr}(T^a T^b) = \delta^{ab}/2 \). \( G_\mu^\nu \) and \( F_\mu^\nu \) are the gluon and photon field strength, respectively. \( O_2 \), \( O_1 \) are the tree level and QCD corrected operators. \( O_3-6 \) are the strong gluon induced operators. \( O_7-10 \) are the electroweak penguin operators due to \( \gamma \) and \( Z \) exchange, and “box” diagrams at loop level. \( O_{11,12} \) are the DPOs.

The Wilson Coefficients (WC) \( c_i \) at a particular scale \( \mu \) are obtained by first calculating the WCs at \( m_W \) scale and then using the renormalization group equation to evolve them to \( \mu \). We have carried out this analysis using the next-to-leading order QCD corrected WCs following Ref [10]. Using \( \alpha_s(m_Z) = 0.118 \), \( \alpha_{em}(m_Z) = 1/128 \), \( m_t = 176 \text{ GeV} \) and \( \mu \approx m_b = 5 \text{ GeV} \), we obtain

\[
\begin{align*}
c_1 &= -0.3125, \\
c_2 &= 1.1502, \\
c_3 &= 0.0174, \\
c_4 &= -0.0373, \\
c_5 &= 0.0104, \\
c_6 &= -0.0459, \\
c_7 &= 1.398 \times 10^{-5}, \\
c_8 &= 3.919 \times 10^{-4}, \\
c_9 &= -0.0103, \\
c_{10} &= 1.987 \times 10^{-3}.
\end{align*}
\]

(2)

The DPO coefficients at the two loop level have the following values [11]:

\[
\begin{align*}
c_{11} &= -0.299, \\
c_{12} &= -0.634.
\end{align*}
\]

(3)

The above Hamiltonian allows one to study the full set of charmless nonleptonic B meson decays. It is clear that only in \( b \to s \) transitions, the DPO may be important because of the enhancement factor \( |V_{tb} V_{ts}^*/V_{ub} V_{us}^*| \approx 57 \) in Eq.(1) compared with the tree contribution. The three processes mentioned before are all of this type. There are two types of DPO contributions. One from \( O_{11} \) and the other from \( O_{12} \). The contribution from \( O_{12} \) is suppressed by a factor of \( \alpha_{em}/\alpha_s \) compared with the contributions from \( O_{11} \). We will therefore neglect its contribution in the following analysis.

The inclusive decay: \( b \to sq\bar{q} \)

Let us first consider \( b \to ss\bar{s} \). This is a pure penguin induced process and does not involve uncertainties associated with hadronic matrix elements. It serves as a clear indication of
the importance of DPO. It also might be feasible to isolate this process experimentally by selecting energetic kaons. In this process the gluon in the DPO has to break up into a pair of \( s\bar{s} \). In our calculation we will use the four-quark operator generated by \( O_{11} \),

\[
H_{11} = i \frac{G_F}{\sqrt{2}} \frac{\alpha_s}{\pi k^2} V_{tb} V_{ts}^* c_{11} m_b \bar{s}\sigma_{\mu\nu} R T^a b \bar{q} \gamma^\mu T^a q k^\nu, \tag{4}
\]

where \( k = p_b - p_s \) in the gluon momentum, and \( q = s \) for \( b \to s\bar{s}s \).

Since there are two s quarks in \( b \to s\bar{s}s \) decay, we have to anti-symmetrize the amplitude. We have

\[
A(b(p_b) \to s_1(p_1)s_2(p_2)s(p_s)) = -\frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* [A_{12} \bar{s}_1 \gamma_\mu L b \bar{s}_2 \gamma^\mu L s + B_{12} \bar{s}_1 \gamma_\mu L b \bar{s}_2 \gamma^\mu R s + C_{12} \bar{s}_1 \gamma_\mu L b \bar{s}_2 \gamma^\mu R s + D_{12} (p_b + p_1)_\mu \bar{s}_1 \gamma_\mu T^a \bar{s}_2 \gamma^\mu T^a s - (1 \leftrightarrow 2)], \tag{5}
\]

where

\[
A_{12} = 4c_3 + 4c_4 - 2c_9 - 2c_{10} - \frac{\alpha_s c_{11} m_b^2}{2\pi} \left( \frac{1}{N_c s} - \frac{1}{t} \right),
\]

\[
B_{12} = 4c_5 - c_7 - \frac{\alpha_s c_{11} m_b^2}{2\pi} \frac{1}{N_c s},
\]

\[
C_{12} = -2(4c_6 - 2c_8 + \frac{\alpha_s c_{11} m_b^2}{2\pi} \frac{1}{t}),
\]

\[
D_{12} = -\frac{\alpha_s c_{11} m_b}{\pi} \frac{1}{s}, \tag{6}
\]

where \( s = (p_b - p_1)^2 \), \( t = (p_b - p_2)^2 \), and \( N_c = 3 \) is the number of colors. The coefficients \((A_{21}, B_{21}, C_{21}, D_{21})\) are obtained by exchanging \( s \) and \( t \) in the above equations. The decay width is given by

\[
\Gamma = \frac{1}{(2\pi)^3} \frac{1}{32m_b^3} \int dsdt |A|^2, \tag{7}
\]

where the integration limits for \( s \) are integrated \( 4m_s^2 \) and \( (m_b - m_s)^2 \), and the limits for \( t \) are \((m_b^2 - s + 3m_s^2 \pm \sqrt{(1 - 4m_s^2/s)((m_b^2 - s - m_s^2)^2 - 4sm_s^2)})/2 \). The integrand \(|A|^2\) is given by
\[ |A|^2 = \frac{1}{2N_c2!} \left( \frac{G_F}{2} \right)^2 \left\{ [4N_c^2|A_{12}|^2 + |A_{21}|^2] + 8N_c Re(A_{12}A_{21}^{*}) \right\} (s + t)(m_b^2 - s - t) \\
+ 2N_c^2 [4|B_{12}|^2 + |C_{21}|^2] t(m_b^2 - t) + (4|B_{21}|^2 + |C_{12}|^2) s(m_b^2 - s)] \\
+ 2(N_c^2 - 1) [2D_{21}^2 t(m_b^2 - t) + D_{21}^2 s(m_b^2 - t)] (m_b^2 - s - t) \\
- 2N_c^2 \left( \frac{1}{N_c} Re(D_{12}D_{21}^{*}) st(m_b^2 - s - t) \\
- 4N_c [Re(B_{12}C_{21}^{*}) t(m_b^2 - t) + Re(C_{12}B_{21}^{*}) s(m_b^2 - s)] \\
+ 2(N_c^2 - 1) m_b [Re(2A_{12}D_{21}^{*} - C_{12}D_{21}^{*}) s \\
+ Re(2D_{12}A_{21}^{*} - D_{12}C_{21}^{*}) t](m_b^2 - s - t) \right\} . \] (8)

Here we have averaged the initial spin and color factors, and also taken into account the symmetrization factor 2!.

In our numerical evaluation, we will use \(|V_{ts}| \approx |V_{cb}| = 0.04, \alpha_s(m_b) = 0.25, m_b = 5\text{GeV}.\)

Since s quarks hadronize into kaons, it seems appropriate to use the constituent mass. We use \(m_s = 0.5\ \text{GeV}.\) Using these numbers, we calculate the relative branching ratio \(R = \Gamma(b \to s\bar{s}s\bar{s})/\Gamma(b \to c\bar{e}\nu_e).\) Without the DPO effects, \(R = 0.018.\) When DPO effects are included, \(R = 0.024.\) We see that the DPO enhances the branching ratio by about 30% and should not be ignored. This result improves upon the result without QCD corrections discussed in Ref. \[12].\)

It is essential to know the distribution of events with respect to invariant masses for selecting the desired events against background. To this end we present two useful spectra. In Figures 1 and 2, we plot the normalized event distributions with respect to diquark invariant mass \(x = (p_1 + p_2)^2/(m_b - m_s)^2\) and quark anti-quark invariant mass \(y = (p_1 + p_s)^2/(m_b - m_s)^2.\) Although the effect of DPO on the overall branching ratio is significant, we see that its effect on the normalized distributions is very small.

We now consider \(b \to sq\bar{q} with q\bar{q} summed over u\bar{u}, d\bar{d} and s\bar{s}.\) This process may be easier to measure experimentally because this process can be isolated by selecting single kaons having high momenta. In this calculation we do not need to antisymmetrize the amplitudes for \(b \to su\bar{u}\) and \(b \to sd\bar{d}.\) We again use constituent masses with \(m_u = m_d = 0.3\ \text{GeV} and \)
$m_s = 0.5 \text{ GeV}$. We find that the effect of DPO becomes weaker. It enhances the branching ratio by about 10% compared with the branching ratio without DPO contribution. The event distribution (shown in Fig 2) with respect to the $q\bar{q}$ invariant mass $y = (p_q + p_{\bar{q}})^2/(m_b - m_s)^2$ is similar to $b \to s\bar{s}$. The DPO effect on the distribution is negligible. The ratio $R = \Gamma(b \to sq\bar{q})/\Gamma(b \to ce\bar{v}_c)$ now also depends on the $b \to u$ tree operators ($O_{1,2}$) which involves the weak phase $\gamma$. We find $R$ is between 0.11 $\sim$ 0.6 when varying the angle $\gamma$ between $0^0 \sim 180^0$.

The semi-inclusive decay: $b \to s\phi$

From experience we know that processes involving hadronic matrix element are difficult to handle because there is no satisfactory method to calculate these matrix elements at present. To have an idea how important the DPO effects are, we use factorization approximation to evaluate the DPO contribution to $b \to s\phi$.

In the factorization approximation, the effective Hamiltonian in Eq.(4) gives vanishing matrix element without Fierz transformation because $s\bar{s}$ operator is in color octet states. Only the Fierz transformed Hamiltonian contributes. After a Fierz transformation, we obtain

$$H_{\text{eff,F}} = \frac{G_F}{\sqrt{2}} \frac{\alpha_s}{4\pi k^2} V_{tb} V_{ts}^* \frac{N_c^2 - 1}{N_c} c_{11} m_b [2m_b \bar{s}^\alpha \gamma_\mu L s^\beta \bar{s}^\gamma \gamma^\mu L b_{b'} - 4m_b \bar{s}^\alpha R s_\beta \bar{s}^\gamma L b_{b'} + (p_{b'}^h + p_{s}^h) (\bar{s}^\alpha \gamma_\mu L s_\beta \bar{s}^\gamma R b_{b'} + \bar{s}^\alpha R s_\beta \bar{s}^\gamma \gamma_\mu L s_\bar{b} ) \gamma^\beta (\delta_\alpha^\beta \delta_\alpha' \gamma^\beta + \frac{2N_c}{N_c^2 - 1} T^\alpha T_{\alpha'})]. \tag{9}$$

The second term in the last bracket gives vanishing contribution in the factorization approximation. We have

$$A(b \to s\phi) = -\frac{G_F}{\sqrt{2}} \frac{\alpha_s}{16\pi k^2} V_{tb} V_{ts}^* \frac{N_c^2 - 1}{N_c} c_{11} m_b i g_\phi \epsilon_\mu [m_b \bar{s}^\gamma \gamma^\mu L b + 4(1 + \frac{m_s^2}{m_\phi^2}) p_{b'}^\mu \bar{s} R b - m_s (1 - \frac{2m_b^2 + 2m_l^2 - m_s}{m_\phi^2}) \bar{s} \gamma^\mu R b - 4 \frac{m_s m_b}{m_\phi^2} p_{b'}^\mu \bar{s} L b] \tag{10},$$

where $\epsilon_\mu$ is the polarization vector of $\phi$. In the above, we have used the parametrization $<\phi|\bar{s}\gamma_\mu s|0> = i g_\phi \epsilon_\mu$, with $g_\phi^2 = 0.0586 \text{ GeV}^4$; and $<\phi|\bar{s}\gamma_{\mu\nu} s|0> = \delta(\epsilon^\mu p_{b'}^\nu - p_{\phi}^\mu \epsilon^\nu)$ with the quark model relation $\delta = -2m_s g_\phi / m_\phi^2$. We assumed that the $s$ and $\bar{s}$ quarks inside $\phi$ each
have similar momentum $p_\phi/2$. In this approximation, $k^2 = (m_b^2 - m_\phi^2/2 + m_s^2)/2$. Using the approximation $A(b \rightarrow s\phi) \approx A(B \rightarrow X_s\phi)$, for $N_c = 3$ we find that without the DPO contribution the branching ratio $BR(B \rightarrow X_s\phi)$ is $1.0 \times 10^{-4}$. With the DPO contribution, $BR(B \rightarrow X_s\phi) = 1.2 \times 10^{-4}$. The branching ratio is enhanced by about 20% when the DPO contribution is included. We also checked the DPO effect with $N_c = 2$ favored by phenomenological fitting of the experimental data \[13\]. In this case the branching ratio is about two times larger, but the DPO contribution only enhances the branching ratio by about 10%.

**The exclusive decay: $B \rightarrow K\pi$**

Here we carry out a calculation for $B^- \rightarrow \pi^- K^0$ as an example. We, again, use the factorization approximation to estimate the amplitude. We have

\[
A_{11}(B^- \rightarrow \pi^- K^0) = -\frac{G_F}{\sqrt{2} 4\pi k^2} V_{tb} V_{ts}^* \frac{N_c^2 - 1}{N_c^2} c_{11} m_b \\
\times \left[ 2m_b < K^0|\bar{s}\gamma_\mu Ld|0 > < \pi^-|\bar{q}\gamma_\mu Lb|B^- > - 4m_b < K^0|\bar{s}Rd|0 > < \pi^-|\bar{d}Lb|B^- > + (p_b + p_s)_\mu (< K^0|\bar{s}\gamma_\mu Ld|0 > < \pi^-|\bar{d}Rb|B^- > + < K^0|\bar{s}Rd|0 > < \pi^-|\bar{d}\gamma_\mu Rb|B^- > - i < K^0|\bar{s}\sigma_{\mu\nu} Rd|0 > < \pi^-|\bar{d}\sigma_{\mu\nu} Rb|B^- > + i < K^0|\bar{s}\gamma_\mu Ld|0 > < \pi^-|\bar{d}\sigma_{\mu\nu} Rb|B^- > \right].
\]

(11)

In the above we have neglected the annihilation contributions of the form $< \pi^- K^0|\bar{s}\Gamma_1 u|0 > < 0|\bar{u}\Gamma_2 |b|B^- >$, where $\Gamma_i$ indicate the appropriate gamma matrices with corresponding Lorentz indices. These contributions are generally much smaller than the ones considered here \[1\].

We proceed with the Lorentz decomposition of the typical hadronic matrix elements:

\[
< K^0|\bar{s}\gamma_\mu \gamma_5 d|0 > = i q^\mu f_K, \text{ with } f_K = 160 \text{ MeV}, \\
< \pi^- \big| d\gamma^\mu (1 \pm \gamma_5) b \big| B^- > = (p_B + p_\pi)\mu f^+ + (p_B - p_\pi)\mu f^-, \\
< \pi^- \big| d\sigma^\mu\nu (1 + \gamma_5) b \big| B^- > = -i 2 h (p_B^\mu p_\pi^\nu - p_B^\nu p_\pi^\mu). \quad (12)
\]
We use the Heavy Quark Effective Theory (HQET) relation \[ \frac{h}{f^+ - f^-} = \frac{1}{4m_b} \] to relate \( h \) to \( f^\pm \).

In terms of the decay constant \( f_K \) and form factors \( f^\pm \), the \( B^- \rightarrow \bar{K}^0\pi^- \) decay amplitude is given by

\[
A_{11}(B^- \rightarrow \bar{K}^0\pi^-) = i \frac{G_F}{\sqrt{2}}\frac{\alpha_s}{4\pi k^2} V_{tb} V_{ts}^* c_{11} \left( \frac{1}{N_c^2} f_K \left[ \left( \frac{m_B^2 - m_\pi^2}{N_c^2} \right) f_{B\pi}^+ + m_K^2 f_{B\pi}^- \right] \right) \times \left( \frac{m_b}{2} + \frac{m_b m_K^2}{m_b - m_d(m_s + m_d)} \right) + \frac{m_K^2}{8(m_b - m_d)} \left( \frac{2m_B^2 - m_K^2}{m_B^2} f_{B\pi}^+ + \frac{1}{2} \left( m_B^2 + 2m_K^2 - m_\pi^2 \right) f_{B\pi}^- \right) \right) \]

\[
+ \frac{1}{8} h_{B\pi} \left( m_B^4 + m_K^4 + m_\pi^4 - 2m_B^2 m_K^2 - 2m_B^2 m_\pi^2 - 2m_K^2 m_\pi^2 \right) \right]. \tag{13}
\]

In the above we have used the same approximation for the kinematics as in the case of \( b \rightarrow s\phi \). We assume that the two quarks s and d inside \( \bar{K}^0 \) share the meson momentum equally, and the b quark carries all the \( B \) meson momentum. In this case \( k^2 = (m_B^2 - m_K^2/2 + m_\pi^2)/2 \)

Using the form factors evaluated in Ref. [13], we obtain the branching ratio \( BR(B^- \rightarrow \pi^-\bar{K}^0) \) to be \( 0.83 \times 10^{-5} \) without the DPO contribution. When the DPO contribution is included, \( BR(B^- \rightarrow \pi^-\bar{K}^0) = 1.0 \times 10^{-5} \). Again, the DPO enhances the branching ratio by about 20%. The dependence on \( N_c \) is weak in this case.

From the above discussions we clearly see that the DPO has important effects on the branching ratio of B decays. These effects enhances the branching ratios by 20% to 30%. All analyses which involve numerical values of the strong penguin contribution are affected. These analyses need to be re-examined. However, since the DPO has the same transformation properties under \( SU(3) \) flavor (and Isospin) symmetry as the ordinary strong penguin operators, many of the analyses related to CP violation in B decays which only depends on transformation properties under flavour symmetry \[ [14][15][16] \] and processes where the strong penguin does not contribute \[ [14] \], will not be affected.
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FIG. 1. The solid and dashed lines are for the event distribution with respect to the normalized diquark invariant mass $x$ with and without the dipole penguin operator contribution, respectively.

FIG. 2. The solid and dashed lines are for the event distribution with respect to normalized quark anti-quark invariant mass $y$ for $b \to s s \bar{s}$ with and without the dipole penguin operator contribution, respectively. The dotted line is for $b \to s q \bar{q}$ with dipole penguin operator contribution.