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Understanding the penguin amplitude in $B \to \phi K$ decays

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Abstract

We calculate branching ratios for pure penguin decay modes, $B \to \phi K$ decays using perturbative QCD approach. Our results of branching ratios are consistent with the experimental data and larger than those obtained from the naive factorization assumption and the QCD-improved factorization approach. This is due to a dynamical penguin enhancement in perturbative QCD approach.

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1. Introduction

Recently the branching ratios of $B \to \phi K$ decays have been measured by the BaBar [1], BELLE [2] and CLEO [3] Collaborations. There is an interesting problem related a penguin contribution to decay amplitudes [4]. A naive estimate of the loop diagram leads to $P/T \sim \alpha_s/(12\pi)\log(m_t^2/m_c^2) \sim \mathrm{O}(0.01)$ where P is a penguin amplitude and T is a tree amplitude. But experimental data for $\mathrm{Br}(B \to K\pi)$ and $\mathrm{Br}(B \to \pi\pi)$ leads to $P/T \sim \mathrm{O}(0.1)$. Therefore, there must be a dynamical enhancement of the penguin amplitude. This problem is studied by Keum et al. using perturbative QCD (PQCD) approach [5]. $B \to \phi K$ modes are important understanding penguin dynamics, because these modes are dominated by penguins. Here we report our study of $B \to \phi K$ decays using PQCD.

PQCD method for inclusive decays was developed by many authors over many years, and this formalism has been successful. Recently, PQCD has been applied to exclusive B meson decays, $B \to K\pi$ [5], $\pi\pi$ [6], $\pi\rho$, $\pi\omega$ [7], KK [8] and $K\eta^{(\prime)}$ [9]. PQCD approach is based on the three scale factorization theorem [10,11]. For example, $B \to K$ transition form factor can be written as

$$F^{BK} \sim \int [dx][db] C_i(t) \Phi_K(x_2, b_2) H(t) \Phi_B(x_1, b_1) \exp\left[-\sum_{j=1,2} \int_{1/b_j}^{t} \frac{d\bar{\mu}}{\bar{\mu}} \gamma_\phi\big(\alpha_s(\bar{\mu})\big)\right], \tag{1}$$

where x_1 and x_2 are momentum fractions of partons, b_1 and b_2 are conjugate variables of parton transverse momenta k_{1T} and k_{2T}, and γ_ϕ is the anomalous dimension of mesons. The hard part $H(t)$ can be calculated

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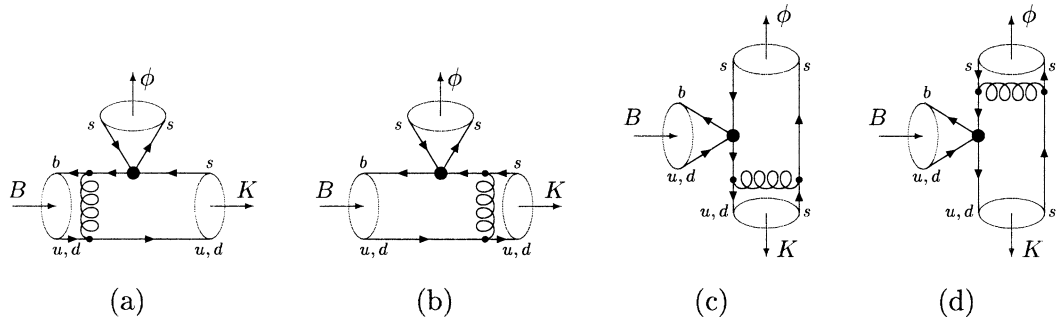

Fig. 1. Feynman diagrams contributing to factorizable amplitudes for $B \to \phi K$.

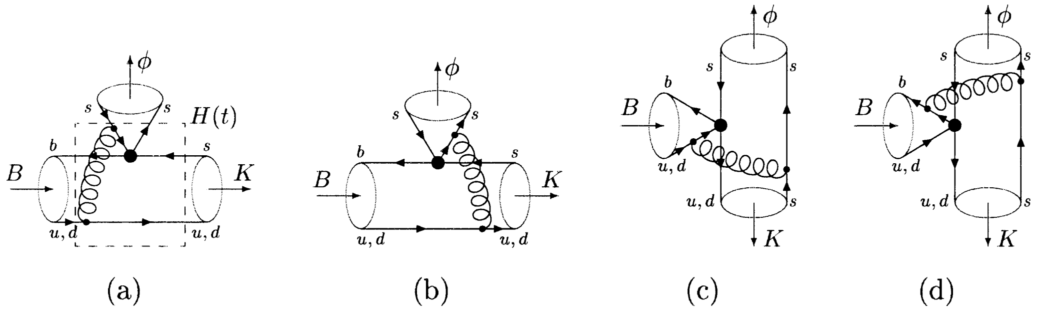

Fig. 2. Feynman diagrams contributing to nonfactorizable amplitudes for $B \to \phi K$.

perturbatively. $C_i(t)$ is the Wilson coefficient corresponding to the four-quark operator causing $B \to K$ transition. The scale t is given explicitly in terms of x_1, x_2, b_1, b_2 and M_B, and it is of $O(\sqrt{\bar{\Lambda} M_B})$. Here $\bar{\Lambda} = M_B - m_b$, where M_B and m_b are B meson mass and b quark mass, respectively. It is important to note that in PQCD, the scale of the Wilson coefficient t can reach below $M_B/2$. In the factorization assumption [12], this scale is fixed at $M_B/2$. The Wilson coefficient for a penguin operator increases as the scale evolves down. This explains the enhancement of the penguin amplitude in PQCD compared to the amplitude obtained by the factorization assumption.

In this method, we can calculate not only factorizable amplitudes but also nonfactorizable and annihilation amplitudes. In case of $B \to \phi K$ decays, the factorizable amplitudes which can be written in terms of form factors F^{BK} and $F^{\phi K}$ are shown in Fig. 1(a)–(d). The nonfactorizable amplitudes are shown in Fig. 2(a)–(d). Ellipses denote meson wave functions in these figures. For illustration purposes, we show the hard part of the nonfactorizable diagram as the dashed box in Fig. 2(a). The parameters in meson wave functions are calculated from the light-cone QCD sum rules, and the theoretical uncertainty of the parameters is about 30%. The hard part depends on the particular processes, but it is calculable. The wave functions contain nonperturbative dynamics and are not calculable, but once it is known, it can be used for other decay processes.

In this Letter, we calculate branching ratios for $B \to \phi K$ modes using PQCD approach. The detail is discussed in Ref. [13]. We predict the branching ratios for $B \to \phi K$ decays, and our predictions agree with the current experimental data and are larger than the values obtained from the naive factorization assumption (FA) and the QCD-improved factorization (QCDF) [14,15].

2. $B \to \phi K$ amplitudes

We consider B meson to be at rest. In the light-cone coordinate, the B meson momentum P_1, the K meson momentum P_2 and the ϕ meson momentum P_3 are taken to be

$$P_1 = \frac{M_B}{\sqrt{2}}(1, 1, \mathbf{0}_T), \qquad P_2 = \frac{M_B}{\sqrt{2}}\left(1 - r_\phi^2, 0, \mathbf{0}_T\right), \qquad P_3 = \frac{M_B}{\sqrt{2}}\left(r_\phi^2, 1, \mathbf{0}_T\right), \tag{2}$$

where $r_\phi = M_\phi / M_B$, and the K meson mass is neglected. The momentum of the spectator quark in the B meson is written as k_1. Since the hard part is independent of k_1^+, the $\delta(k_1^+)$ function appears after integrating over its conjugate spacial variable. Therefore, k_1 has only the minus component k_1^- and small transverse components $\mathbf{k}_{1T}$. k_1^- is given as $k_1^- = x_1 P_1^-$, where x_1 is a momentum fraction. The quarks in the K meson have plus components $x_2 P_2^+$ and $(1 - x_2) P_2^+$, and the small transverse components $\mathbf{k}_{2T}$ and $-\mathbf{k}_{2T}$, respectively. The quarks in the ϕ meson have the minus components $x_3 P_3^-$ and $(1 - x_3) P_3^-$, and the small transverse components $\mathbf{k}_{3T}$ and $-\mathbf{k}_{3T}$, respectively. The ϕ meson longitudinal polarization vector ϵ_ϕ and two transverse polarization vector $\epsilon_{\phi T}$ are given by $\epsilon_\phi = (1/\sqrt{2}\, r_\phi)(-r_\phi^2, 1, \mathbf{0}_T)$ and $\epsilon_{\phi T} = (0, 0, \mathbf{1}_T)$.

The B meson wave function for incoming state and the K and ϕ meson wave functions for outgoing state with up to twist-3 terms are written as

$$\Phi^{(\mathrm{in})}_{B,\alpha\beta,ij} = \frac{i\delta_{ij}}{\sqrt{2N_c}} \int dx_1\, d^2\mathbf{k}_{1T}\, e^{-i(x_1 P_1^- z_1^+ - \mathbf{k}_{1T}\mathbf{z}_{1T})} \left[(\not{P}_1 + M_B)\gamma_5 \phi_B(x_1, \mathbf{k}_{1T})\right]_{\alpha\beta}, \tag{3}$$

$$\Phi^{(\mathrm{out})}_{K,\alpha\beta,ij} = \frac{-i\delta_{ij}}{\sqrt{2N_c}} \int_0^1 dx_2\, e^{ix_2 P_2 \cdot z_2} \gamma_5 \left[\not{P}_2 \phi_K^A(x_2) + m_{0K}\phi_K^P(x_2) + m_{0K}\left(\not{v}\not{n} - 1\right)\phi_K^T(x_2)\right]_{\alpha\beta}, \tag{4}$$

$$\Phi^{(\mathrm{out})}_{\phi,\alpha\beta,ij} = \frac{\delta_{ij}}{\sqrt{2N_c}} \int_0^1 dx_3\, e^{ix_3 P_3 \cdot z_3} \left[M_\phi \not{\epsilon}_\phi \phi_\phi(x_3) + \not{\epsilon}_\phi \not{P}_3 \phi_\phi^t(x_3) + M_\phi \phi_\phi^s(x_3)\right]_{\alpha\beta}, \tag{5}$$

where i and j is color indices, and α and β are Dirac indices. m_{0K} is related to the chiral symmetry breaking scale, $m_{0K} = M_K^2/(m_d + m_s)$. v and n are defined as $v^\mu = P_2^\mu / P_2^+$ and $n^\mu = z_2^\mu / z_2^- = (0, 1, \mathbf{0}_T)$. We neglect the terms which are proportional to the transverse polarization vector ϵ_ϕ^T, because these terms drop out from our calculation kinematically. The explicit form of these wave functions will be shown in Section 3.

Widths of $B \to \phi K$ decays can be expressed as

$$\Gamma = \frac{G_F^2}{32\pi M_B} |\mathcal{A}|^2. \tag{6}$$

The decay amplitudes, $\mathcal{A}$, and $\bar{\mathcal{A}}$, corresponding to $B^0 \to \phi K^0$, and $\overline{B}^0 \to \phi \overline{K}^0$, respectively, are written as

$$\mathcal{A} = f_\phi V_{ts} V_{tb}^* F_e^P + V_{ts} V_{tb}^* \mathcal{M}_e^P + f_B V_{ts} V_{tb}^* F_a^P + V_{ts} V_{tb}^* \mathcal{M}_a^P, \tag{7}$$

$$\overline{\mathcal{A}} = f_\phi V_{ts}^* V_{tb} F_e^P + V_{ts}^* V_{tb} \mathcal{M}_e^P + f_B V_{ts}^* V_{tb} F_a^P + V_{ts}^* V_{tb} \mathcal{M}_a^P, \tag{8}$$

where F_e is the amplitude for factorizable diagrams which are considered in FA. F_a and $\mathcal{M}$ are the annihilation factorizable and the nonfactorizable diagrams which are neglected in FA. The indices e and a, denote the tree topology, and annihilation topology, respectively. The index P denotes the contribution from diagrams with a penguin operator. The decay amplitudes $\mathcal{A}^+$ and $\mathcal{A}^-$, corresponding to $B^+ \to \phi K^+$, and $B^- \to \phi K^-$, respectively, are written as

$$\mathcal{A}^+ = f_\phi V_{ts} V_{tb}^* F_e^P + V_{ts} V_{tb}^* \mathcal{M}_e^P + f_B V_{ts} V_{tb}^* F_a^P + V_{ts} V_{tb}^* \mathcal{M}_a^P - f_B V_{us} V_{ub}^* F_a^T - V_{us} V_{ub}^* \mathcal{M}_a^T, \tag{9}$$

$$\mathcal{A}^- = f_\phi V_{ts}^* V_{tb} F_e^P + V_{ts}^* V_{tb} \mathcal{M}_e^P + f_B V_{ts}^* V_{tb} F_a^P + V_{ts}^* V_{tb} \mathcal{M}_a^P - f_B V_{us}^* V_{ub} F_a^T - V_{us}^* V_{ub} \mathcal{M}_a^T, \tag{10}$$

where the index T denotes tree contributions. Since the Cabibbo–Kobayashi–Maskawa (CKM) matrix elements for the tree amplitudes are much smaller than for the penguin amplitudes, the tree contributions are very small.

The factorizable diagrams are given as Fig. 1. The factorizable penguin amplitude, F_e^P, which comes from Fig. 1(a) and Fig. 1(b) is written as

$$\begin{aligned} F_e^P &= 8\pi C_F M_B^4 \int_0^1 dx_1\, dx_2 \int_0^\infty b_1 db_1 b_2 db_2 \phi_B(x_1, b_1) \\ &\quad \times \Big\{ \big[(1+x_2)\phi_K^A(x_2) + r_K(1-2x_2)\big(\phi_K^P(x_2) + \phi_K^T(x_2)\big)\big] E_e\big(t_e^{(1)}\big) N_t \big\{x_2(1-x_2)\big\}^c h_e(x_1, x_2, b_1, b_2) \\ &\quad + 2r_K \phi_K^P(x_2) E_e\big(t_e^{(2)}\big) N_t \big\{x_1(1-x_1)\big\}^c h_e(x_2, x_1, b_2, b_1) \Big\}, \end{aligned} \tag{11}$$

where $N_t\{x(1-x)\}^c$ is the factor for the threshold resummation [16]. We use $N_t = 1.775$ and $c = 0.3$ [17]. The evolution factors are defined by $E_e(t) = \alpha_s(t) a_e(t) \exp[-S_B(t) - S_K(t)]$ where $\exp[-S_i(t)]$ is the factor for the k_T resummation [18,19]. The explicit forms of the factor $S_i(t)$ are given, for example, in Ref. [5]. The hard scales $t_e^{(1)}$ and $t_e^{(2)}$, which are the scales in hard process, are given by $t_e^{(1)} = \max(\sqrt{x_2}\, M_B, 1/b_1, 1/b_2)$ and $t_e^{(2)} = \max(\sqrt{x_1}\, M_B, 1/b_1, 1/b_2)$. The Wilson coefficient is given by

$$a_e(t) = C_3 + \frac{C_4}{N_c} + C_4 + \frac{C_3}{N_c} + C_5 + \frac{C_6}{N_c} - \frac{1}{2}\left(C_7 + \frac{C_8}{N_c} + C_9 + \frac{C_{10}}{N_c} + C_{10} + \frac{C_9}{N_c}\right). \tag{12}$$

The hard function, which is the Fourier transformation of the virtual quark propagator and the hard gluon propagator, is given by

$$h_e(x_1, x_2, b_1, b_2) = K_0\big(\sqrt{x_1 x_2}\, M_B b_1\big)\big[\theta(b_1 - b_2) K_0\big(\sqrt{x_2}\, M_B b_1\big) I_0\big(\sqrt{x_2}\, M_B b_2\big) + (b_1 \leftrightarrow b_2)\big]. \tag{13}$$

The factorizable annihilation diagrams shown in Fig. 1(c) and Fig. 1(d), and the nonfactorizable diagrams shown in Fig. 2(a)–(d) can be also calculated in the same way as F_e^P [13].

3. Numerical results

We use the model of the B meson wave function written as

$$\phi_B(x, b) = N_B x^2 (1-x)^2 \exp\left[-\frac{1}{2}\left(\frac{x M_B}{\omega_B}\right)^2 - \frac{\omega_B^2 b^2}{2}\right], \tag{14}$$

where $\omega_B = 0.40$ GeV [20]. N_B is determined by normalization condition given by

$$\int_0^1 dx\, \phi_B(x, b = 0) = \frac{f_B}{2\sqrt{2N_c}}. \tag{15}$$

The K meson wave functions are given as

$$\phi_K^A(x) = \frac{f_K}{2\sqrt{2N_c}} 6x(1-x)\left[1 + 3a_1(1-2x) + \frac{3}{2}a_2\big\{5(1-2x)^2 - 1\big\}\right], \tag{16}$$

$$\begin{aligned} \phi_K^P(x) &= \frac{f_K}{2\sqrt{2N_c}}\left[1 + \frac{1}{2}\left(30\eta_3 - \frac{5}{2}\rho_K^2\right)\big\{3(1-2x)^2 - 1\big\}\right. \\ &\quad \left. - \frac{1}{8}\left(3\eta_3\omega_3 + \frac{27}{20}\rho_K^2 + \frac{81}{10}\rho_K^2 a_2\right)\big\{3 - 30(1-2x)^2 + 35(1-2x)^4\big\}\right], \end{aligned} \tag{17}$$

$$\phi_K^T(x)=\frac{f_K}{2\sqrt{2N_c}}(1-2x)\left[1+6\left(5\eta_3-\frac{1}{2}\eta_3\omega_3-\frac{7}{20}\rho_K^2-\frac{3}{5}\rho_K^2a_2\right)\left(1-10x+10x^2\right)\right], \tag{18}$$

where $\rho_K=(m_d+m_s)/M_K$ [21,22]. The parameters of these wave functions are given as $a_1=0.17$, $a_2=0.20$, $\eta_3=0.015$ and $\omega_3=-3.0$ where the renormalization scale is 1 GeV.

The ϕ meson wave functions are given as

$$\phi_\phi(x)=\frac{f_\phi}{2\sqrt{2N_c}}6x(1-x), \tag{19}$$

$$\phi_\phi^t(x)=\frac{f_\phi^T}{2\sqrt{2N_c}}\left[3(1-2x)^2+\frac{35}{4}\zeta_3^T\left\{3-30(1-2x)^2+35(1-2x)^4\right\}\right.$$
$$\left.+\frac{3}{2}\delta_+\left\{1-(1-2x)\log\frac{1-x}{x}\right\}\right], \tag{20}$$

$$\phi_\phi^s(x)=\frac{f_\phi^T}{4\sqrt{2N_c}}\left[(1-2x)\left(6+9\delta_++140\zeta_3^T-1400\zeta_3^Tx+1400\zeta_3^Tx^2\right)+3\delta_+\log\frac{x}{1-x}\right], \tag{21}$$

where $\zeta_3^T=0.024$ and $\delta_+=0.46$ [23]. We have found that the final results are insensitive to the values chosen for ζ_3^T and δ_+.

We use the Wolfenstein parameters for the CKM matrix elements $A=0.819$, $\lambda=0.2196$, $R_b\equiv\sqrt{\rho^2+\eta^2}=0.38$ [24], and choose the angle $\phi_3=\pi/2$ [5]. We have found that the final results are quite insensitive to the values of ϕ_3. For the values of meson masses, we use $M_B=5.28$ GeV, $M_K=0.49$ GeV and $M_\phi=1.02$ GeV. In addition, for the values of meson decay constants, we use $f_B=190$ MeV, $f_K=160$ MeV, $f_\phi=237$ MeV and $f_\phi^T=215$ MeV. The B meson life times are given as $\tau_{B^0}=1.55\times10^{-12}$ s and $\tau_{B^\pm}=1.65\times10^{-12}$ s. And we use $\Lambda_{\mathrm{QCD}}^{(4)}=0.250$ GeV and $m_{0K}=1.70$ GeV [5].

We show the numerical results of each amplitude for $B^0\to\phi K^0$ and $B^\pm\to\phi K^\pm$ decays in Table 1. The factorizable penguin amplitude F_e^P gives a dominant contribution to $B\to\phi K$ decays. The factorizable annihilation penguin amplitude F_a^P generates a large strong phase. In $B^\pm\to\phi K^\pm$ modes, there are contributions from $f_BF_a^T$ and M_a^T. These tree amplitudes contribute only a few percent to the whole amplitude, since the CKM matrix elements related to the tree amplitudes are very small. In order to isolate the trivial uncertainty from f_B, f_K and f_ϕ, we express our prediction for $B\to\phi K$ as

$$\mathrm{Br}\left(B^0\to\phi K^0\right)=\left|\frac{f_Bf_Kf_\phi}{190\ \mathrm{MeV}\ 160\ \mathrm{MeV}\ 237\ \mathrm{MeV}}\right|^2\times\left(9.43\times10^{-6}\right), \tag{22}$$

$$\mathrm{Br}\left(B^\pm\to\phi K^\pm\right)=\left|\frac{f_Bf_Kf_\phi}{190\ \mathrm{MeV}\ 160\ \mathrm{MeV}\ 237\ \mathrm{MeV}}\right|^2\times\left(10.1\times10^{-6}\right). \tag{23}$$

We found that our result is insensitive to f_ϕ^T/f_ϕ. For example, 10% variation of f_ϕ^T/f_ϕ leads to less than 1% variation in our final result. The current experimental values are summarized in Table 2. The values which are predicted in PQCD are consistent with the current experimental data. However, our branching ratios have the theoretical error from the $\mathrm{O}(\alpha_s^2)$ corrections, the higher twist corrections, and the error of input parameters. Large uncertainties come from the meson decay constants, the shape parameter ω_B, and m_{0K}. These parameters are fixed from the other modes ($B\to K\pi, D\pi$, etc.). We try to vary ω_B from 0.36 to 0.44 GeV, then we obtain $\mathrm{Br}(B^\pm\to\phi K^\pm)=(7.54\sim13.9)\times10^{-6}$. Next, we set $\omega_B=0.40$ and try to vary m_{0K} from 1.40 to 1.80 GeV, then we obtain $\mathrm{Br}(B^\pm\to\phi K^\pm)=(6.65\sim11.4)\times10^{-6}$.

Now, we consider the ratio of the branching ratio for the $B^0\to\phi K^0$ decay to the one for the $B^+\to\phi K^+$ decay. The theoretical uncertainties from various parameters are small, since the parameters in the numerator cancel out those in the denominator. The difference between the two branching ratios come in principle from B

Table 1
Contribution to $B^0 \to \phi K^0$ and $B^\pm \to \phi K^\pm$ decays from each amplitude

	$B^0 \to \phi K^0$	$B^\pm \to \phi K^\pm$
$f_\phi F_e^P$	-1.03×10^{-1}	-1.03×10^{-1}
$f_B F_a^P$	$6.45 \times 10^{-3} + i4.28 \times 10^{-2}$	$6.17 \times 10^{-3} + i4.20 \times 10^{-2}$
M_e^P	$5.24 \times 10^{-3} - i3.61 \times 10^{-3}$	$5.24 \times 10^{-3} - i3.61 \times 10^{-3}$
M_a^P	$-8.03 \times 10^{-4} - i1.73 \times 10^{-3}$	$-6.56 \times 10^{-4} - i7.22 \times 10^{-4}$
$f_B F_a^T$		$-1.11 \times 10^{-1} - i3.75 \times 10^{-2}$
M_a^T		$1.60 \times 10^{-2} + i2.77 \times 10^{-2}$

Table 2
The experimental data of $B \to \phi K$ branching ratios from BaBar [1], BELLE [2] and CLEO [3]

	$\mathrm{Br}(B^0 \to \phi K^0)$	$\mathrm{Br}(B^\pm \to \phi K^\pm)$
BaBar	$(8.1^{+3.1}_{-2.5} \pm 0.8) \times 10^{-6}$	$(7.7^{+1.6}_{-1.4} \pm 0.8) \times 10^{-6}$
BELLE	$(8.7^{+3.8}_{-3.0} \pm 1.5) \times 10^{-6}$	$(10.6^{+2.1}_{-1.9} \pm 2.2) \times 10^{-6}$
CLEO	$< 12.3 \times 10^{-6}$	$(5.5^{+2.1}_{-1.8} \pm 0.6) \times 10^{-6}$

meson life times, tree and electroweak penguin contributions in annihilation amplitudes. We found that the tree and electroweak penguin amplitudes in the annihilation diagrams are negligible. Tree amplitudes are suppressed by two factors. First, they are annihilation processes which are suppressed by helicity. Second, they are multiplied by small CKM matrix elements. We predict that this ratio is

$$\frac{\mathrm{Br}(B^0 \to \phi K^0)}{\mathrm{Br}(B^+ \to \phi K^+)} = 0.95, \tag{24}$$

where the theoretical uncertainties from m_{0K} and ω_B are less than 1%. The ratio is essentially given by the life time difference. The experimental value of this ratio from BELLE [2] is $\mathrm{Br}(B^0 \to \phi K^0)/\mathrm{Br}(B^+ \to \phi K^+) = 0.82^{+0.39}_{-0.32} \pm 0.10$.

In FA, the branching ratio is very sensitive to the effective number of colors N_c^{eff}. If we set $N_c^{\mathrm{eff}} = 3$, then the branching ratio is about 4.5×10^{-6} where the scale of the Wilson coefficient is taken to $M_B/2$ and F^{BK} is 0.38 from the BSW model. In QCDF, branching ratios for $B \to \phi K$ decays are predicted as $\mathrm{Br}(B^0 \to \phi K^0) = (4.0^{+2.9}_{-1.4}) \times 10^{-6}$ and $\mathrm{Br}(B^- \to \phi K^-) = (4.3^{+3.0}_{-1.4}) \times 10^{-6}$ including the annihilation diagram [15]. Our predicted values are larger than these results. This is due to the enhancement of the Wilson coefficient for the penguin amplitude as explained in Section 1. In PQCD approach, the scale of the Wilson coefficients, which is equal to the hard scale t, can reach lower values than $M_B/2$.

4. Summary

In this Letter, we calculate $B^0 \to \phi K^0$ and $B^\pm \to \phi K^\pm$ decays in PQCD approach. Our predicted branching ratios agree with the current experimental data and are larger than the values obtained by FA and QCDF. Because the Wilson coefficients for penguin operators are enhanced dynamically in PQCD.

Note added

After this work has been completed, we become aware of a similar calculation by Chen et al. [25]. Our results are in agreement.

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References

[1] G. Cavoto (for BaBar Collaboration), hep-ex/0105018.
[2] A. Bozek (for BELLE Collaboration), Talk presented at the International Europhysics Conference on High Energy Physics, Budapest, Hungary, July 12–18, 2001.
[3] CLEO Collaboration, Phys. Rev. Lett. 86 (2001) 3718.
[4] For example, see, I.I. Bigi, A.I. Sanda, CP Violation, Cambridge, 2000.
[5] Y.Y. Keum, H.-N. Li, A.I. Sanda, Phys. Lett. B 504 (2001) 6;
Y.Y. Keum, H.-N. Li, A.I. Sanda, Phys. Rev. D 63 (2001) 054008.
[6] C.D. Lü, K. Ukai, M.Z. Yang, Phys. Rev. D 63 (2001) 074009.
[7] C.D. Lü, M.Z. Yang, hep-ph/0011238.
[8] C.H. Chen, H.-N. Li, Phys. Rev. D 63 (2001) 014003.
[9] E. Kou, A.I. Sanda, hep-ph/0106159.
[10] C.H. Chang, H.-N. Li, Phys. Rev. D 55 (1997) 5577.
[11] T.W. Yeh, H.-N. Li, Phys. Rev. D 56 (1997) 1615.
[12] M. Bauer, B. Stech, M. Wirbel, Z. Phys. C 34 (1987) 103;
M. Bauer, B. Stech, M. Wirbel, Z. Phys. C 29 (1985) 637.
[13] S. Mishima, hep-ph/0107163.
[14] X.G. He, J.P. Ma, C.Y. Wu, Phys. Rev. D 63 (2001) 094004.
[15] H.Y. Cheng, K.C. Yang, Phys. Rev. D 64 (2001) 074004.
[16] H.-N. Li, hep-ph/0102013.
[17] T. Kurimoto, H.-N. Li, A.I. Sanda, hep-ph/0105003.
[18] J.C. Collins, D.E. Soper, Nucl. Phys. B 193 (1981) 381.
[19] J. Botts, G. Sterman, Nucl. Phys. B 325 (1989) 62.
[20] M. Bauer, M. Wirbel, Z. Phys. C 42 (1989) 671.
[21] P. Ball, JHEP 09 (1998) 005.
[22] P. Ball, JHEP 01 (1999) 010.
[23] P. Ball, V.M. Braun, Y. Koike, K. Tanaka, Nucl. Phys. B 529 (1998) 323.
[24] Particle Data Group, Eur. Phys. J. C 15 (2000) 1.
[25] C.H. Chen, Y.Y. Keum, H.-N. Li, hep-ph/0107165.