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$B \to f_0(980)\,K$ DECAYS IN QCD FACTORIZATION

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Dedicated to the memory of Professor Dubravko Tadić

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The decay $B \to f_0(980)K$ is studied within the framework of QCD factorization. Its decay rate is suppressed relative to $B \to \pi^0 K$ owing to a destructive interference between $(S-P)(S+P)$ and $(V-A)(V-A)$ penguin contributions. The interference between the $(S-P)(S+P)$ penguin contributions arising from the strange and light quark components of $f_0(980)$ is destructive for $\pi/2 > \theta > 0$ and constructive for $-\pi/2 < \sin\theta < 0$, with θ being the mixing angle of strange and nonstrange quark contents of $f_0(980)$ in the two-quark picture for light scalar mesons. A negative mixing angle, as preferred by several $f_0(980)$ production experiments, is also supported by the measurement of $B \to f_0(980)K$ decay. We conclude that the short-distance interactions are not adequate to explain the experimental observation of $f_0(980)K^+ > \pi^0 K^+$ and $f_0(980)K^0 \gtrsim \pi^0 K^0$ decay rates. Possible mechanisms for the enhancement of $f_0(980)K$ are discussed.

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1. Introduction

The decay of the B meson into a scalar meson $f_0(980)$ was first measured by Belle [1] in the charged B decays to $K^\pm\pi^\mp\pi^\pm$ and a large branching fraction product for the $f_0(980)K^\pm$ final states was found. A recent updated result by Belle yields [2]

$$\mathcal{B}(B^+ \to f_0(980)K^+ \to \pi^+\pi^-K^+) = (7.55 \pm 1.24^{+1.63}_{-1.18}) \times 10^{-6}. \quad (1)$$

The Belle result is subsequently confirmed by the BaBar measurement [3]:

$$\mathcal{B}(B^+ \to f_0(980)K^+ \to \pi^+\pi^-K^+) = (9.2 \pm 1.2^{+2.1}_{-2.6}) \times 10^{-6}. \quad (2)$$

The weighted average is [4]

$$\mathcal{B}(B^+ \to f_0(980)K^+ \to \pi^+\pi^-K^+) = (8.49^{+1.35}_{-1.26}) \times 10^{-6}. \quad (3)$$

BaBar has also measured the neutral mode $B^0 \to f_0(980)K^0$ with the result [5]

$$\mathcal{B}(B^0 \to f_0(980)K^0 \to \pi^+\pi^-K^0) = (6.0 \pm 0.9 \pm 1.3) \times 10^{-6}. \quad (4)$$

This channel is of special interest as possible indications of physics beyond the Standard Model (SM) may be observed in the time-dependent CP asymmetries in the penguin-dominated B decays such as $B^0 \to f_0(980)K_S$. . The mixing-induced CP-violation parameter S is expected to be $-\sin\beta$ in the SM. The most recent measurements by BaBar and Belle yield

$$\sin\beta(f_0K_S) = \begin{cases} 0.95^{+0.23}_{-0.32} \pm 0.10 & \text{BaBar F [6]} \\ -0.47 \pm 0.41 \pm 0.08 & \text{Belle [7]}\,. \end{cases} \quad (5)$$

The deviation from $\sin 2\beta = 0.726 \pm 0.037$ [8] derived from the decay $B \to J/\psi K_S$ may hint at a possible new physics.

In order to extract the absolute branching ratios for $B \to f_0K$, we use the value of $\Gamma(f_0 \to \pi\pi)/[\Gamma(f_0 \to \pi\pi) + \Gamma(f_0 \to K\overline{K})] \approx 0.68$ [9] to obtain $\mathcal{B}(f_0(980) \to \pi^+\pi^-) \approx 0.45$ and

$$\begin{aligned} \mathcal{B}(B^+ \to f_0(980)K^+) &\approx (18.9^{+3.0}_{-2.8}) \times 10^{-6}, \\ \mathcal{B}(B^0 \to f_0(980)K^0) &\approx (13.3 \pm 3.6) \times 10^{-6}. \end{aligned} \quad (6)$$

Comparing with the averaged branching ratios, $(12.1 \pm 0.8) \times 10^{-6}$ for $B^+ \to \pi^0K^+$ and $(11.5 \pm 1.0) \times 10^{-6}$ for $B^0 \to \pi^0K^0$ [4], we see that for the decay rates $f_0(980)K^+ > \pi^0K^+$ and $f_0(980)K^0 \gtrsim \pi^0K^0$.

This decay mode has been studied in Refs. [10] and [11] within the framework of the pQCD approach based on the k_{T} factorization theorem. It is found that the branching ratio is of order 5×10^{-6} (see Fig. 2 of Ref. [11]), which is smaller than the measured value by a factor of $3 \sim 4$. In the present paper, we wish to re-examine this decay and see if the discrepancy between theory and experiment can be resolved in the QCD factorization approach [12, 13, 14].

2. $B \to f_0(980)K$ decays in QCD factorization

2.1. Framework

To proceed, we first discuss the decay constants and form factors. The decay constants are defined by

$$\langle K(p)|A_\mu|0\rangle = -\mathrm{i}f_K p_\mu, \qquad \langle f_0|V_\mu|0\rangle = 0, \qquad \langle f_0|\bar{q}q|0\rangle = m_{f_0}\tilde{f}_q\,. \quad (7)$$

The scalar meson $f_0(980)$ cannot be produced via the vector current owing to charge conjugation invariance or conservation of vector current. The decay constant $\tilde{f}_q$ will be discussed later. Form factors for $B \to P$ and $B \to S$ transitions (P: pseudoscalar meson, S: scalar meson) are defined by [15]

$$\langle P(p_P)|V_\mu|B(p_B)\rangle = \left(p_{B\mu}+p_{P\mu}-\frac{m_B^2-m_P^2}{q^2}\,q_\mu\right) F_1^{BP}(q^2)+\frac{m_B^2-m_P^2}{q^2}q_\mu\,F_0^{BP}(q^2), \tag{8}$$

where $q_\mu = (p_B - p_P)_\mu$, and [16][1]

$$\langle S(p_S)|A_\mu|B(p_B)\rangle = -\mathrm{i}\left[\left(p_{B\mu}+p_{S\mu}-\frac{m_B^2-m_S^2}{q^2}\,q_\mu\right) F_1^{BS}(q^2) + \frac{m_B^2-m_S^2}{q^2}q_\mu\,F_0^{BS}(q^2)\right]. \tag{9}$$

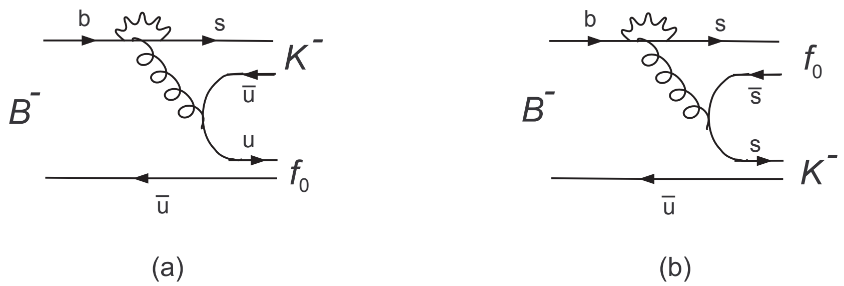

Fig. 1. Penguin contributions to $B^- \to f_0(980)K^-$.

The penguin-dominated $B^- \to f_0K^-$ receive two different types of penguin contributions as depicted in Fig. 1. Within the framework of QCD factorization [12], its decay amplitude reads

$$\begin{aligned} A(B^- \to f_0K^-) =& \\ &-\frac{G_F}{\sqrt{2}}\left\{\lambda_u\,[a_1+a_4^u+a_{10}^u-2(a_6^u+a_8^u)r_\chi]+\lambda_c\,[a_4^c+a_{10}^c-2(a_6^c+a_8^c)r_\chi]\right\} \\ &\times f_K(m_B^2-m_{f_0}^2)F_0^{Bf_0^u}(m_K^2) \qquad (10) \\ &-\left\{\lambda_u(2a_6'^u-a_8'^u)+\lambda_c(2a_6'^c-a_8'^c)\right\}\tilde{f}_s\,\frac{m_{f_0}}{m_b}\,(m_B^2-m_K^2)F_0^{BK}(m_{f_0}^2) \\ &+\mathcal{A}_{ann}(B^- \to f_0K^-), \end{aligned}$$

[1] As shown in Ref. [16], a factor of $(-i)$ is needed in Eq. (9) in order for the $B \to S$ form factors to be positive. This also can be checked from heavy quark symmetry [16].

where $\lambda_q \equiv V_{qb}V_{qs}^*$, and $\mathcal{A}_{ann}$ is the weak annihilation contribution

$$\mathcal{A}_{\rm ann}(B^- \to f_0 K^-) = \frac{G_F}{\sqrt{2}}\Bigg\{\lambda_u\, c_2 \mathcal{A}_1^i + (\lambda_u+\lambda_c)\Bigg[(c_3+c_9)\mathcal{A}_1^i + (c_5+c_7)\mathcal{A}_i^3 + N_c\Big[c_6+c_8+\frac{1}{N_c}(c_5+c_7)\Big]\mathcal{A}_3^f\Bigg]\Bigg\}, \tag{11}$$

where $\mathcal{A}_3^f$ is the factorizable annihilation amplitude induced from $(S-P)(S+P)$ operator and $\mathcal{A}_{1,3}^i$ are nonfactorizable ones induced from $(V-A)(V-A)$ and $(S-P)(S+P)$ operators, respectively. The explicit expressions of $\mathcal{A}_{1,3}^i$ and $\mathcal{A}_3^f$ are given by (see also Refs. [13, 14])

$$\mathcal{A}_1^i = \kappa\int_0^1 {\rm d}x{\rm d}y\left\{\Phi_{f_0}(x)\Phi_K(y)\left[\frac{1}{y(1-x\bar{y})}+\frac{1}{\bar{x}^2y}\right] + \frac{4\mu_\chi m_{f_0}}{m_b^2}\Phi_{f_0}^p(x)\Phi_K^p(y)\,\frac{2}{\bar{x}y}\right\},$$

$$\mathcal{A}_3^i = \kappa\int_0^1 {\rm d}x{\rm d}y\left\{\frac{2\mu_\chi}{m_b}\Phi_{f_0}(x)\Phi_K^p(y)\,\frac{2\bar{y}}{\bar{x}y(1-x\bar{y})} - \frac{2m_{f_0}}{m_b}\Phi_K(y)\Phi_{f_0}^p(x)\,\frac{2x}{\bar{x}y(1-x\bar{y})}\right\},$$

$$\mathcal{A}_3^f = \kappa\int_0^1 {\rm d}x{\rm d}y\left\{\frac{2\mu_\chi}{m_b}\Phi_{f_0}(x)\Phi_K^p(y)\,\frac{2(1+\bar{x})}{\bar{x}^2y} + \frac{2m_{f_0}}{m_b}\Phi_K(y)\Phi_{f_0}^p(x)\,\frac{2(1+y)}{\bar{x}y^2}\right\}, \tag{12}$$

where $\kappa \equiv (C_F/N_c^2)\pi\alpha_s f_B f_K(\tilde{f}_s - \tilde{f}_u)$ with $C_F = (N_c^2-1)/(2N_c)$, and Φ_M (Φ_M^p) is the twist-2 (twist-3) light-cone distribution amplitude of the meson M.

In Eq. (10), $r_\chi(\mu) = m_K^2/[m_b(\mu)(m_u(\mu)+m_s(\mu))]$ and the expressions for the parameters a_i^q ($q=u,c$) will be discussed shortly. The superscript u of the form factor $F_0^{Bf_0^u}$ reminds us that it is the u quark component of f_0 involved in the form factor transition [see Fig. 1(a)]. In contrast, the subscript s of the decay constant $\tilde{f}_s$ indicates that it is the strange quark component responsible for the penguin contribution of Fig. 1(b).

For comparison, we also write down the $B^- \to \pi^0 K^-$ decay amplitude [17]

$$\begin{aligned} &A(B^- \to \pi^0 K^-) \\ &= {\rm i}\frac{G_F}{2}\Bigg\{\lambda_u\left[a_1+a_4^u+a_{10}^u+2(a_6^u+a_8^u)r_\chi\right]+\lambda_c\left[a_4^c+a_{10}^c+2(a_6^c+a_8^c)r_\chi\right]\Bigg\} \\ &\times f_K(m_B^2-m_\pi^2)F_0^{B\pi}(m_K^2) \\ &+\frac{\rm i}{\sqrt{2}}\Big[\lambda_u a_2 + \frac{3}{2}(\lambda_u+\lambda_c)(-a_7+a_9)\Big]f_\pi(m_B^2-m_K^2)F_0^{BK}(m_\pi^2) \\ &+{\rm i}\mathcal{A}_{\rm ann}(B^-\to\pi^0K^-). \end{aligned} \tag{13}$$

We see that a_4 and a_6 terms contribute constructively to $\pi^0 K^-$ but destructively to $f_0 K^-$ decay.

The parameters a_i^q with $q = u, c$ can be calculated in the QCD factorization approach [12]. They are basically the Wilson coefficients in conjunction with short-distance nonfactorizable corrections such as vertex corrections and hard spectator interactions. In general, they have the expressions [13, 14]

$$a_i^q(M_1 M_2) = c_i + \frac{c_{i\pm1}}{N_c} + \frac{c_{i\pm1}}{N_c}\frac{C_F\alpha_s}{4\pi}\Big[V_i(M_2) + \frac{4\pi^2}{N_c}H_i(M_1M_2)\Big] + P_i^q(M_2), \quad (14)$$

where $i = 1, \cdots, 10$, the upper (lower) signs apply when i is odd (even), M_1 is the emitted meson and M_2 shares the same spectator quark with the B meson. The quantities $V_i(M_2)$ account for vertex corrections, $H_i(M_1M_2)$ for hard spectator interactions with a hard gluon exchange between the emitted meson and the spectator quark of the B meson and $P_i(M_2)$ for penguin contractions. The explicit expressions of these quantities can be found in [13, 14], in particular, Eq. (46) of Ref. [13], except that the hard spectator function $H_{K\pi}$ is replaced by H_{Kf_0} which reads

$$H_{Kf_0} = \frac{\tilde{f}_u f_B}{F_0^{Bf_0^u}(0) m_B^2} \int_0^1 \frac{\mathrm{d}\rho}{\rho}\,\Phi_B(\rho) \int_0^1 \frac{\mathrm{d}\xi}{\bar{\xi}}\,\Phi_K(\xi) \int_0^1 \frac{\mathrm{d}\eta}{\bar{\eta}} \left[\Phi_{f_0}(\eta) + \frac{2m_{f_0}}{m_b}\,\frac{\bar{\xi}}{\xi}\,\Phi_{f_0}^p(\eta)\right], \quad (15)$$

where $\bar{\xi} \equiv 1 - \xi$. As for the parameters $a_{6,8}'^q$ appearing in Eq. (10), they have the same expressions as $a_{6,8}^q$ except that the function G_K (see Eq. (50) of Ref. [13]) is replaced by G_{f_0}, Φ_K by Φ_{f_0}, $\hat{G}_K$ (see Eq. (55) of Ref. [13]) by $\hat{G}_{f_0}$ and Φ_K^p by $\Phi_{f_0}^p$. Formally, $a_i(i \neq 6, 8)$ and $a_{6,8} r_\chi$ should be renormalization scale and scheme independent. In practice, there exists some residual scale dependence in $a_i(\mu)$ to finite order.

2.2. *Distribution amplitudes*

In the present paper we will take the asymptotic forms for kaon twist-2 and twist-3 distribution amplitudes:

$$\Phi_K(x) = 6x(1 - x), \qquad \Phi_K^p(x) = 1. \quad (16)$$

As for the distribution amplitude of $f_0(980)$, it needs some elaboration.

It is known that the underlying structure of scalar mesons is not well established theoretically (for a review, see e.g. Refs. [18, 19, 20]). It has been suggested that the light scalars below or near 1 GeV – the isoscalars $f_0(600)$ (or σ), $f_0(980)$, the isodoublet κ and the isovector $a_0(980)$ – form an SU(3) flavor nonet, while scalar mesons above 1 GeV, namely, $f_0(1370)$, $a_0(1450)$, $K_0^*(1430)$ and $f_0(1500)/f_0(1710)$, form another nonet. A consistent picture [20] provided by the data suggests that the scalar meson states above 1 GeV can be identified as a conventional $q\bar{q}$ nonet with some possible glue content, whereas the light scalar mesons below or near 1 GeV form predominately a $qq\bar{q}\bar{q}$ nonet [21, 22] with a possible mixing with 0^+ $q\bar{q}$ and glueball states. This is understandable because in the $q\bar{q}$ quark model, the 0^+ meson has a unit of orbital angular momentum and hence it should have a higher mass above 1 GeV. On the contrary, four quarks $q^2\bar{q}^2$ can form a 0^+ meson without introducing a unit of orbital angular momentum. Moreover, color and spin-dependent interactions favor a flavor nonet configuration with attraction between the qq and $\bar{q}\bar{q}$ pairs. Therefore, the 0^+ $q^2\bar{q}^2$ nonet has a mass near or below 1 GeV. This four-quark scenario explains naturally the mass degeneracy of $f_0(980)$ and $a_0(980)$, the broader decay widths of $\sigma(600)$ and $\kappa(800)$ than $f_0(980)$ and $a_0(980)$, and the large coupling of $f_0(980)$ and $a_0(980)$ to $K\bar{K}$.

While the above-mentioned four-quark assignment of $f_0(980)$ is certainly plausible when the light scalar meson is produced in low-energy reactions, it is dubious that the energetic $f_0(980)$ produced in B decays is dominated by the four-quark configuration as it requires to pick up two energetic quark-antiquark pairs to form a fast-moving light four-quark scalar meson. The Fock states of $f_0(980)$ consists of $q\bar{q}$, $q^2\bar{q}^2$, $q\bar{q}g$ etc. Naively, it is thus expected that the distribution amplitude Φ_{f_0} would be smaller in the four-quark model than in the two-quark picture. Then one will not be able to explain the observed $B \to f_0(980)K$ decays.

In the naive 2-quark picture, $f_0(980)$ is purely an $s\bar{s}$ state and this is supported by the data of $D_s^+ \to f_0\pi^+$ and $\phi \to f_0\gamma$ implying the copious $f_0(980)$ production via its $s\bar{s}$ component. However, there also exists some experimental evidence indicating that $f_0(980)$ is not purely an $s\bar{s}$ state. First, the observation of $\Gamma(J/\psi \to f_0\omega) \approx \frac{1}{2}\Gamma(J/\psi \to f_0\phi)$ [23] clearly indicates the existence of the non-strange and strange quark content in $f_0(980)$. Second, the fact that $f_0(980)$ and $a_0(980)$ have similar widths and that the f_0 width is dominated by $\pi\pi$ also suggests the composition of $u\bar{u}$ and $d\bar{d}$ pairs in $f_0(980)$; that is, $f_0(980) \to \pi\pi$ should not be OZI suppressed relative to $a_0(980) \to \pi\eta$. Therefore, isoscalars $\sigma(600)$ and f_0 must have a mixing

$$|f_0(980)\rangle = |s\bar{s}\rangle\cos\theta + |n\bar{n}\rangle\sin\theta, \qquad |\sigma_0(500)\rangle = -|s\bar{s}\rangle\sin\theta + |n\bar{n}\rangle\cos\theta, \quad (17)$$

with $n\bar{n} \equiv (\bar{u}u + \bar{d}d)/\sqrt{2}$. The distribution amplitudes Φ_s and Φ_n corresponding to $f_0^s = \bar{s}s$ and $f_0^n = \bar{n}n \equiv (\bar{u}u + \bar{d}d)/\sqrt{2}$, respectively, are

$$\begin{aligned}
\langle f_0^n(p)|\bar{q}(z)\gamma_\mu q(0)|0\rangle &= p_\mu \tilde{f}_n \int_0^1 \mathrm{d}x\, e^{\mathrm{i}xp\cdot z}\Phi_n(x), \\
\langle f_0^s(p)|\bar{s}(z)\gamma_\mu s(0)|0\rangle &= p_\mu \tilde{f}_s \int_0^1 \mathrm{d}x\, e^{\mathrm{i}xp\cdot z}\Phi_s(x), \\
\langle f_0^n(p)|\bar{n}(z)n(0)|0\rangle &= m_{f_0} \tilde{f}_n \int_0^1 \mathrm{d}x\, e^{\mathrm{i}xp\cdot z}\Phi_n^p(x),
\end{aligned}$$

$$\langle f_0^s(p)|\bar{s}(z)s(0)|0\rangle \;=\; m_{f_0}\tilde{f}_s\int_0^1 \mathrm{d}x\, e^{\mathrm{i}xp\cdot z}\Phi_s^p(x) \tag{18}$$

where $\tilde{f}_q$ is defined in Eq. (7). They satisfy the relations $\Phi_{n,s}(x) = -\Phi_{n,s}(1-x)$ due to charge-conjugation invariance (that is, the distribution amplitude vanishes at $x=1/2$) and $\Phi^p_{n,s}(x) = \Phi^p_{n,s}(1-x)$ and hence $\int_0^1 \mathrm{d}x\,\Phi_{n,s}(x)=0$ and $\int_0^1 \mathrm{d}x\,\Phi^p_{n,s}(x)=1$. For the scalar meson made of $q\bar{q}$, its general distribution amplitude has the form [24]

$$\Phi_S(x) = 6x(1-x)\left[B_0 + \sum_{n=1}^{\infty} B_n\, C_n^{3/2}(1-2x)\right], \tag{19}$$

where B_0, B_n are constants and $C_n^{3/2}$ is the Gegenbauer polynomial. For the isosinglet scalar mesons σ and f_0, $B_0 = 0$. Hence, the leading twist-2 distribution amplitude for f_0 reads

$$\Phi_{f_0}(x) = 6B_1\, x(1-x)(3-6x). \tag{20}$$

In the present work, we shall use $B_1 = 1.1$ as inferred from the analysis in Ref. [24]. As for the twist-3 distribution amplitude $\Phi^p_{f_0}(x)$, its asymptotic form is the same as the light pseudoscalar meson to the leading conformal expansion [25]. Hence, we take

$$\Phi^p_{f_0}(x) = 1. \tag{21}$$

In the $q\bar{q}$ description of $f_0(980)$, it follows from that

$$F_0^{B^- f_0} = \frac{1}{\sqrt{2}}\sin\theta\, F_0^{B^- f_0^{u\bar{u}}}, \qquad F_0^{B^0 f_0} = \frac{1}{\sqrt{2}}\sin\theta\, F_0^{B^0 f_0^{d\bar{d}}}, \tag{22}$$

where the superscript $q\bar{q}$ denotes the quark content of f_0 involved in the transition. The form factor for B to the scalar meson transition has been calculated in the covariant light-front model [16]. From Table VI of Ref. [16], it is clear that $F_0^{Bf_0^{q\bar{q}}}(0)$ with $q\bar{q} = u\bar{u}$ or $d\bar{d}$ is of order 0.25 which is very similar to $F_0^{B\pi}(0)$. Based on the sum-rule technique, the decay constant f_s defined by $\langle f_0^s|\bar{s}s|0\rangle = m_{f_0}f_s$ has been estimated in Refs. [26] and [27] with similar results, namely, $f_s \approx 0.18$ GeV. However, this quantity is scale-dependent. For our purpose, we need to evolute it from the typical sum rule scale of the order of 0.5 GeV to $\mu = 2.1$ GeV. It turns out that $f_s(2.1\,\mathrm{GeV}) \approx 0.30$ GeV [32]. In the two-quark scenario, the decay constants $\tilde{f}_s$ and $\tilde{f}_u$ are related to f_s by

$$\tilde{f}_s = f_s\cos\theta, \qquad \tilde{f}_u = f_s\sin\theta/\sqrt{2}. \tag{23}$$

Experimental implications for the $f_0-\sigma$ mixing angle have been discussed in detail in Ref. [28]. A typical mixing angle is $\theta \approx \pm 35°$. As pointed out in Ref. [28], the solution $\theta \sim -35°$ is preferred by the measurements of $J/\psi \to f_0\phi$ and $J/\psi \to f_0\omega$, the $f_0(980)$ coupling to $\pi\pi$ and $K\bar{K}$ and the radiative decays $\phi \to f_0\gamma$ and $f_0 \to \gamma\gamma$. As we shall see shortly, a negative f_0–σ mixing angle is also supported by the measurement of $B \to f_0(980)K$ decays.

In the four-quark picture, $f_0(980)$ has the flavor function $s\bar{s}(u\bar{u}+d\bar{d})/\sqrt{2}$. However, the estimate of its decay constant and form factors is beyond the conventional quark model.

Using the asymptotic distribution amplitudes of the kaon and $f_0(980)$, the annihilation contributions are simplified to

$$\begin{aligned}
\mathcal{A}_1^i &\approx \kappa\left[18B_1(3\pi^2-10)+\frac{8\mu_\chi m_{f_0}}{m_b^2}X_A^2\right], \\
\mathcal{A}_3^i &\approx 12\kappa\left[\frac{3\mu_\chi}{m_b}B_1X_A(-X_A+4)-\frac{m_{f_0}}{m_b}X_A(3X_A-2)\right], \\
\mathcal{A}_3^f &\approx 12\kappa\left[-\frac{\mu_\chi}{m_b}B_1(6X_A-11)+\frac{m_{f_0}}{m_b}X_A(2X_A-1)\right],
\end{aligned} \tag{24}$$

where the endpoint divergence $X_A \equiv \int_0^1 \mathrm{d}x/x$ is parametrized as [13]

$$X_A = \ln\left(\frac{m_B}{\Lambda_h}\right)(1+\rho_A e^{\mathrm{i}\phi_A}) \tag{25}$$

with Λ_h being a hadronic scale of order 500 MeV and ρ_A a real parameter $0 \le \rho_A \le 1$.

3. Results and Discussion

It is ready to perform numerical calculations. At the scale $\mu = 2.1$ GeV, the numerical results for the relevant a_i^q are

$$\begin{aligned}
&a_4^u = -0.0366 - \mathrm{i}\,0.0137, && a_4^c = -0.0423 - \mathrm{i}\,0.0054, \\
&a_6^u = -0.0583 - \mathrm{i}\,0.0122, && a_6^c = -0.0616 - \mathrm{i}\,0.0034, \\
&a_8^u = (74.0 - \mathrm{i}\,4.5)\times 10^{-5}, && a_8^c = (73.2 - \mathrm{i}\,2.4)\times 10^{-5}, \\
&a_{10}^u = (-60.7 + \mathrm{i}\,66.4)\times 10^{-5}, && a_{10}^c = (-62.1 + \mathrm{i}\,68.4)\times 10^{-5}, \\
&a_1 = 1.0739 + \mathrm{i}\,0.0216, \qquad a'^u_{6,8} = a^u_{6,8}, && a'^c_{6,8} = a^c_{6,8}.
\end{aligned} \tag{26}$$

For current quark masses, we use $m_b(m_b) = 4.4$ GeV, $m_c(m_b) = 1.3$ GeV, $m_s(2.1\,\mathrm{GeV}) = 90$ MeV and $m_q/m_s = 0.044$.

In Fig. 2 is shown the branching ratio of $B^- \to f_0(980)K^-$ versus the strange-nonstrange mixing angle θ. It turns out that the annihilation contribution is rather small. When $\theta = 0$, f_0 is a pure $s\bar{s}$ state and hence the penguin diagram Fig. 1(a) does not contribute (i.e. the form factor $F_0^{Bf_0^u}$ vanishes). On the other extreme with $\theta = \pm 90°$, f_0 is purely a $n\bar{n}$ state and the penguin diagram Fig. 1(b) vanishes (i.e. $\tilde{f}_s = 0$). For a finite mixing angle, the interference between a_6^q and $a_6^{\prime q}$ penguin terms arising from Figs. 1(a) and 1(b), respectively, is destructive for $\pi/2 > \theta > 0$ and constructive for $-\pi/2 < \sin\theta < 0$. As stated before, a negative mixing angle is preferred by experiments. It is evident from Fig. 2 that the negative angle solution is also supported by the measurement of $B \to f_0 K$. We obtain $\mathcal{B}(B^- \to f_0 K^-) = 2.8 \times 10^{-6}$ for $\theta = 35°$ and 8.4×10^{-6} for $\theta = -35°$. However, even the maximal branching ratio 8.8×10^{-6} occurring at $\theta \approx -25°$ is still too small by a factor of 2 compared to experiment.

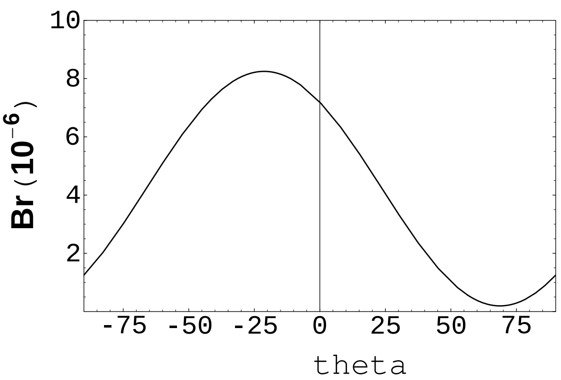

Fig. 2. Branching ratio of $B^- \to f_0(980)K^-$ versus the mixing angle θ of strange and nonstrange components of $f_0(980)$.

The fact that the observed $f_0(980)K^-$ rate is significantly higher than the naive model prediction calls for some mechanisms beyond the conventional short-distance model considerations. Some possibilities are:

- Final state interactions. The predicted $B \to \pi K$ rates in the short-distance approach are in general smaller than the data by around 20% (see e.g. Ref. [29]). Long-distance rescattering via charm intermediate states (or the so-called charming penguins) will not only enhance πK rates but also drive sizable direct CP violation observed recently in the $B^0 \to K^+\pi^-$ mode [29]. The same rescattering effects are expected to enhance $f_0(980)K$ rates by (20 − 30)%.
- Gluonic coupling of the scalar meson. It is known that a possible explanation of the enormous production of $B \to \eta' K$ and $B \to \eta' X_s$ may be ascribed to the process $b \to s + g + g$ and the two gluons fragment into η'. The same mechanism may be also responsible for the enhancement of $f_0(980)K$ [30].

- Subleading corrections arising from the three-parton Fock states of final-state mesons. It has been shown that this effect alone can enhance the branching ratio of $K\eta'$ to the level above 50×10^{-6} [31]. By the same token, it is expected that the three-parton Fock state contributions will play an eminent role for the enhancement of $f_0(980)K$, which we will report in a separate work [32].

4. *Conclusions*

We have studied the decay $B \to f_0(980)K$ using the QCD factorization approach. Its decay rate is suppressed relative to $B \to \pi^0 K$ owing to a destructive interference between a_4 and a_6 penguin contributions. In order to enhance $f_0(980)K$ rates, the interference between the $(S-P)(S+P)$ penguin contributions arising from the strange and light quark components of $f_0(980)$ should be constructive, implying a negative strange–nonstrange mixing angle in the two-quark picture for $f_0(980)$. We conclude that the short-distance interactions are not adequate to explain the observed large $f_0(980)K$ branching ratios. Several possible mechanisms for the enhancement of $f_0(980)K$ are discussed.

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RASPADI B $\rightarrow f_0(980)$ K U QCD FAKTORIZACIJI

Proučava se raspad B $\rightarrow f_0(980)$ K u okviru QCD faktorizacije. Vjerojatnost raspada je potisnuta u odnosu na B $\rightarrow \pi^0$ K raspad zbog destruktivne interferencije pingvinskih doprinosa (S−P)(S+P) i (V−A)(V−A). Interferencija pingvinskih doprinosa (S−P)(S+P), koja nastaje zbog komponenata stranog i laganog kvarka u $f_0(980)$, je destruktivna za $\pi/2 > \theta > 0$ i konstruktivna za $-\pi/2 < \sin\theta < 0$, gdje je θ kut miješanja sadržaja stranog i nestranog kvarka u $f_0(980)$ u dvokvarkovskoj slici lakih skalarnih mezona. Negativan kut miješanja izvodi se u analizama više mjerenja tvorbe $f_0(980)$, uz potvrde mjerenjima B $\rightarrow f_0(980)$ K. Zaključujemo da kratkodosežna međudjelovanja nisu dostatna za objašnjenje eksperimentalnih opažanja da je $f_0(980)K^+ > \pi^0 K^+$ i $f_0(980)K^0 \gtrsim \pi^0 K^0$. Raspravljaju se mogući mehanizmi povećane vjerojatnosti $f_0(980)$K.