$B \to f_0(980) K^{(*)}$ decays and final state interactions

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We study the exclusive decays of $B\to f_0(980)K^{(*)}$ in the framework of perturbative QCD by identifying $f_0(980)$ as the composition of $s\bar{s}$ and $n\bar{n}=(u\bar{u}+d\bar{d})/\sqrt{2}$. We find that the influence of the $n\bar{n}$ content on the predicted branching ratios is crucial. We discuss the possible rescattering and gluonium states which could enhance the branching ratios of the considered decays. We point out that the *CP* asymmetry in $B \to f_0(980)K_{S,L}$ could be a new explorer of $\sin 2\phi_1$.

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Recently, Belle [1] and BaBar [2] have opened new channels in three-body nonleptonic B decays, such as $B \to KKK$, $KK\pi$, and $K\pi\pi$ decays. In particular, Belle has observed the decay of $B^+\to f_0(980)K^+$ with the branching ratio (BR) product of $\mathrm{Br}(B^+\to f_0(980)K^+)\times \mathrm{Br}(f_0(980) \to \pi^+\pi^-)=(9.6^{+2.5+1.5+3.4}_{-2.3-1.5-0.8})\times 10^{-6}$ [1]. Since $f_0(980)$ is a neutral scalar meson, the measured $B\to f_0(980)K$ decays not only show for the first time that B decays to scalar-pseudoscalar final states but also provide the chance to understand the content of $f_0(980)$ and its production.

The essential inner structure of the scalar meson $f_0(980)$ is still obscure since it was established first by Ref. [3] with a phase shift analysis. In the literature, $f_0(980)$ could be four-quark states denoted by $qq\bar{q}\bar{q}$ [4] or $K\bar{K}$ molecular states [5] or $\bar{q}q$ states [6]. However, one objection against the possibility of $K\bar{K}$ states is that the binding energy of 10–20 MeV for $K\bar{K}$ is much smaller than the measured width in the range 40–100 MeV [7]. It is suggested that, in terms of the measured $\phi\to f_0(980)\gamma$ and $f_0(980)\to\gamma\gamma$ [7–10] decays and $D_s^+\to f_0(980)\pi^+$ decay [10,11], the flavor contents of $f_0(980)$ are mostly $s\bar{s}$ and a small portion of $n\bar{n}=(u\bar{u}+d\bar{d})/\sqrt{2}$. In this paper, we take $f_0(980)$ to be composed of $\bar{q}q$ states and use $|f_0(980)\rangle=\cos\phi_s|s\bar{s}\rangle+\sin\phi_s|n\bar{n}\rangle$ to denote its flavor wave function.

Before presenting our perturbative QCD (PQCD) calculation to the decays, we would like to give a brief model-independent analysis of $B\to f_0(980)K$. For simplicity, our analysis will only concentrate on the dominant factorizable parts and regard the $f_0(980)$ as the composition of $s\bar{s}$ so that at the quark level the process corresponds to $b\to s\bar{s}s$ decay. Since, at the quark level, the physics for $B\to\phi K$ decays is the same as $B\to f_0(980)K^{(*)}$, we first examine the ϕK mode. We start by writing the effective Hamiltonian for the $b\to s$ transition as

$$H_{\mathrm{eff}}=\frac{G_F}{\sqrt{2}}\sum_{q'=u,c}V_{q'}\left[C_1(\mu)\mathcal{O}_1^{(q')}+C_2(\mu)\mathcal{O}_2^{(q')}+\sum_{i=3}^{10}C_i(\mu)\mathcal{O}_i\right],$$

where $V_{q'}=V^*_{q's}V_{q'b}$ are the products of the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements, $C_i(\mu)$ are the Wilson coefficients (WCs), and $\mathcal{O}_i$ correspond to the four-quark operators. The explicit expressions of $C_i(\mu)$ and $\mathcal{O}_i$ can be found in Ref. [12]. It is known that the vector meson ϕ has the current matrix elements

$$\langle 0|\bar{s}\gamma_\mu s|\phi\rangle=m_\phi f_\phi\varepsilon_\mu, \quad \langle 0|\bar{s}s|\phi\rangle=0. \tag{1}$$

Therefore the decay amplitude of $B_d\to\phi K^0$ can be simply described by [13]

$$A(B_d\to\phi K^0)=f_\phi V_t^*(a_3^{(s)}+a_4^{(s)}+a_5^{(s)})F_e^{BK}+2f_BV_t^*a_6^{(d)}F_{a6}^{\phi K}+\cdots, \tag{2}$$

where $a_i^{(q)}$ are defined by

$$\begin{aligned}
&a_1=C_1+\frac{C_2}{N_c},\quad a_2=C_2+\frac{C_1}{N_c},\\
&a_{3,4}^{(q)}=C_{3,4}+\frac{3e_q}{2}C_{9,10}+a_{3,4}^{\prime(q)},\\
&a_{3,4}^{\prime(q)}=\frac{C_{4,3}}{N_c}+\frac{3e_q}{2N_c}C_{10,9},\\
&a_{5,6}^{(q)}=C_{5,6}+\frac{3e_q}{2}C_{7,8}+a_{5,6}^{\prime(q)},\\
&a_{5,6}^{\prime(q)}=\frac{C_{6,5}}{N_c}+\frac{3e_q}{2N_c}C_{8,7}.
\end{aligned} \tag{3}$$

Note that a_2 is larger than a_1 and $a_{4,6}^{(q)}$ are much larger than $a_{3,5}^{(q)}$, F_e^{BK} is the $B\to K$ decay form factor and $F_{a6}^{\phi K}$ denotes $\langle\phi K|\bar{s}\gamma_5 d|0\rangle$ annihilation effect. We neglect $\langle\phi K|\bar{s}\gamma_\mu\gamma_5 d|0\rangle$ due to the chirality suppression. According to the results of Ref. [14], the predicted BR of $B_d\to\phi K^0$ is around $\sim 9.6\times 10^{-6}$ which is consistent with Belle's and BaBar's results given by $(10.0^{+1.9+0.9}_{-1.7-1.3})\times 10^{-6}$ [15] and $(8.7^{+1.7}_{-1.5}\pm 0.9)\times 10^{-6}$ [16], respectively. However, in the $B \to f_0(980)K^0$ decay the relevant current matrix elements are given by

$$\langle 0|\bar{s}\gamma_\mu s|f_0(980)\rangle=0, \quad \langle 0|\bar{s}s|f_0(980)\rangle=m_{f_0}\tilde{f}, \tag{4}$$

where $m_{f_0}(\tilde{f})$ are the mass (decay constant) of $f_0(980)$. From Eqs. (1) and (4), it is clear that the situation in the $f_0(980)K^0$ mode is just opposite to the case of ϕK^0, i.e., the role in the vector and scalar vertex is exchanged with each other. Hence the decay amplitude for $B\to f_0(980)K^0$ can be simply written as

$$A(B_d\to f_0(980)K^0)=2r_f\tilde{f}V_t^* a_6^{(s)}S_e^{BK} + 2f_B V_t^* a_6^{(d)}F_{a6}^{f_0K}+\cdots, \qquad (5)$$

where $r_f=m_{f_0}/M_B$, $S_e^{BK}=\langle K|\bar{b}s|B\rangle/M_B$ and the factor of 2 comes from Fierz transformation in the $(V-A)\times(V+A)$ operator. By taking $M_B\approx m_b$ with m_b being the b-quark mass, one can expect $S_e^{BK}\approx F_e^{BK}$ by equation of motion. Comparing Eq. (2) with Eq. (5), it is obvious that there is a suppression factor $r_f=0.186$ in $B\to f_0(980)K^0$. Including the factor of 2, taking $a_6^{(s)}(\sqrt{\bar{\Lambda}M_B})/a_4^{(s)}(\sqrt{\bar{\Lambda}M_B})\approx 1.5$ and $\tilde{f}\sim f_\phi$ and neglecting the annihilation contributions, we estimate the ratio of $\mathrm{Br}(B\to f_0(980)K^0)/\mathrm{Br}(B\to\phi K^0)$ being around 0.31. With the average value of Belle and BaBar, the predicted $\mathrm{Br}(B\to f_0(980)K^0)$ is around 2.8×10^{-6}. Will the value change while we include the $\bar{n}n$ content and annihilation contributions? In the following we display the results in a more serious theoretical approach.

It is known that large uncertainities are always involved in the calculations of transition matrix elements while studying exclusive hadron decays. Nevertheless, the problem will become mild in the heavy B meson decays because of the enormous B samples produced by B factories. That is, more precise measurements will help the theory to pin down the unknown parameters to make some predictions. In this paper, we adopt the PQCD approach in which the applications to exclusive heavy B meson decays, such as $B\to K\pi$ [19], $B\to\pi\pi(KK)$ [20–22], $B\to\phi\pi(K)$ [14,23], $B\to\eta^{(\prime)}K$ [24], and $B\to\rho K$ [13] decays, have been studied and found that all of them are consistent with the current experimental data [25,26]. In the PQCD, in order to solve the various divergences encountered at end point, we will include not only k_T resummation, for removing end-point singularities, but also threshold resummation, for smearing the double logarithmic divergence arisen from weak corrections [17].

In order to satisfy the local current matrix elements with Eq. (4), the $f_0(980)$ meson distribution amplitude is given by

$$\langle 0|\bar{q}(0)_j q(z)_l|f_0\rangle=\frac{1}{\sqrt{2N_c}}\int_0^1 dx e^{-ixP\cdot z} m_f[\mathbf{1}]_{lj}\Phi_{f_0}(x) \qquad (6)$$

with the normalization

$$\int_0^1 dx\Phi_{f_0}(x)=\frac{\tilde{f}}{2\sqrt{2N_c}}.$$

For B and $K^{(*)}$ mesons, the corresponding distribution amplitudes [14,17] are written as

$$\langle 0|\bar{b}(z)_j q(0)_l|B\rangle=\frac{1}{\sqrt{2N_c}}\int_0^1 dx e^{-izxP\cdot x}([\not{P}]_{jl} + M_B[I]_{jl})\gamma_5\phi_B(x) \qquad (7)$$

$$\langle K|\bar{q}(z)_j s(0)_l|0\rangle=\frac{1}{\sqrt{2N_c}}\int_0^1 dx e^{-izP\cdot x} \times\{[\gamma_5\not{P}]_{jl}\Phi_K(x) + m_K^0[\gamma_5]_{jl}\Phi_K^p(x) + m_K^0[\gamma_5(\not{n}_-\not{n}_+-1)]_{jl}\Phi_K^\sigma(x)\}, \qquad (8)$$

$$\langle K^*(\varepsilon_L)|\bar{q}(z)_j s(0)_l|0\rangle=\frac{1}{\sqrt{2N_c}}\int_0^1 dx e^{-izP\cdot x} \times\{M_{K^*}[\not{\varepsilon}_L]_{lj}\phi_{K^*}(x) + [\not{\varepsilon}_L\not{P}]_{lj}\phi_{K^*}^t(x) + M_{K^*}[I]_{lj}\phi_{K^*}^s(x)\} \qquad (9)$$

with $n_-=(0,1,0_\perp)$, $n_+=(1,0,0_\perp)$, and m_K^0 being the chiral symmetry breaking parameter, $\Phi_{K(*)}(x)$ stand for twist-2 wave functions and the remains belong to twist-3 with their explicit expressions being given in Refs. [14,18]. Although K^* meson has three polarizations, only the longitudinal part is involved in the decays of $B\to f_0(980)K^*$. We note that as the K meson, the distribution amplitude of $f_0(980)$ can have more complicated spin structures. However, since our purpose is just for the properties of $B\to f_0(980)K^{(*)}$, we consider only the simplest case. Moreover, if $f_0(980)$ consists of $\bar{s}s$ mostly, the choice of Eq. (6) is clearly dominant.

The decay rates of $B\to f_0(980)K$ are expressed by

$$\Gamma=\frac{G_F^2 M_B^3}{32\pi}|A|^2, \qquad (10)$$

where A includes all possible components of $f_0(980)$ and topologies. As mentioned before, $f_0(980)$ has the components of $\bar{s}s$ and $\bar{n}n$, for different contents, the amplitudes of $B_d\to f_0(980)\bar{K}^0$ and $B^+\to f_0(980)K^+$ are written as

$$A_{\bar{s}s}=\tilde{f}V_t^* S_{e6}^{P(s)}+f_B V_t^* S_{a46}^{P(d)}+\cdots,$$
$$A_{\bar{n}n}=f_K V_t^* N_{e46}^{P(d)}+f_B V_t^* N_{a46}^{P(d)}+\cdots,$$
$$A_{\bar{s}s}^{\pm}=\tilde{f}V_t^* S_{e6}^{P(s)}+f_B V_t^* S_{a46}^{P(u)}-f_B V_u^* S_a+\cdots,$$
$$A_{\bar{n}n}^{\pm}=f_K V_t^* N_{e46}^{P(u)}+f_B V_t^* N_{a46}^{P(u)}-f_K V_u^* N_e-f_B V_u^* N_a+\cdots, \qquad (11)$$

respectively, where $S_{e(a)}^{P(q)}$ ($N_{e(a)}^{P(q)}$) denote the emission (annihilation) contributions of the $\bar{s}s(\bar{n}n)$ content from penguin diagrams while $S_{e(a)}(N_{e(a)})$ are from tree contributions. The total decay amplitude for the neutral (charged) mode is described by $A=\cos\phi A_{ss}^{(+)}+\sin\phi A_{nn}^{(+)}/\sqrt{2}$. Because of the smallness of nonfactorizable effects, we neglect to show them in Eq. (11) and just display the factorizable contributions with emission and annihilation topologies. But we will include their effects in our final numerical results. According to Eqs. (6)–(8), the hard amplitudes for $\bar{s}s$ content are derived as

$$S_{e6}^{P(q)}=16\pi r_f C_F M_B^2\int_0^1 dx_1 dx_3\int_0^\infty b_1 db_1 b_3 db_3 \times\Phi_B(x_1,b_1)\{[\Phi_K(x_3)+2r_K\Phi_K^p(x_3)+r_K x_3[\Phi_K^p(x_3) -\Phi_K^\sigma(x_3)]]ES_{e6}^{(q)}(t_e^{(1)})h_e(x_1,x_3,b_1,b_3) +2r_K\Phi_K^p(x_3)ES_{e6}^{(q)}(t_e^{(2)})h_e(x_3,x_1,b_3,b_1)\}, \quad (12)$$

$$S_{a4}^{P(q)}=-16\pi r_f r_K C_F M_B^2\int_0^1 dx_2 dx_3\int_0^\infty b_2 db_2 b_3 db_3\Phi_{f_0}(x_2) \times\{[(1+x_3)\Phi_K^p(1-x_3)+(1-x_3)\Phi_K^\sigma(1-x_3)] \times E_{a4}^{(q)}(t_a^{(1)})h_a(x_2,x_3,b_2,b_3)-[(1+x_2)\Phi_K^p(1-x_3)] \times E_{a4}^{(q)}(t_a^{(2)})h_a(x_3,x_2,b_3,b_2)\}, \quad (13)$$

$$S_{a6}^{P(q)}=16\pi r_f C_F M_B^2\int_0^1 dx_2 dx_3\int_0^\infty b_2 db_2 b_3 db_3\Phi_{f_0}(x_2) \times\Phi_K(1-x_3)\{2E_{a6}^{(q)}(t_a^{(1)})h_a(x_2,x_3,b_2,b_3) +x_2E_{a6}^{(q)}(t_a^{(2)})h_a(x_3,x_2,b_3,b_2)\} \quad (14)$$

and the results for the content of $\bar{n}n$ are given by

$$N_{e4}^{P(q)}=8\pi r_f C_F M_B^2\int_0^1 dx_1 dx_2\int_0^\infty b_1 db_1 b_2 db_2 \times\Phi_B(x_1,b_1)\Phi_{f_0}(x_2)\{(1-2x_2)EN_{e4}^{(q)} \times(t_e^{(1)})h_e(x_1,x_2,b_1,b_2) +2EN_{e4}^{(q)}(t_e^{(2)})h_e(x_2,x_1,b_2,b_1)\}, \quad (15)$$

$$N_{e6}^{P(q)}=-16\pi r_f r_K C_F M_B^2\int_0^1 dx_1 dx_2 \times\int_0^\infty b_1 db_1 b_2 db_2\Phi_B(x_1,b_1)\Phi_{f_0}(x_2), \times\{(2+x_2)EN_{e6}^{(q)}(t_e^{(1)})h_e(x_1,x_2,b_1,b_2) +2EN_{e6}^{(q)}(t_e^{(2)})h_e(x_2,x_1,b_2,b_1)\}. \quad (16)$$

Here, the hard part functions $h_{e(a)}$, mainly arising from the propagators of hard gluon and valence quark, have included the threshold resummation factor [14], and the evolution factors are given by

$$ES_{e6}^{(q)}(t)=\alpha_s(t)a_6^{(q)}(t)\exp[-S_B(t,x_1)-S_K(t,x_3)],$$

$$EN_{ei}^{(q)}(t)=\alpha_s(t)a_i^{(q)}(t)\exp[-S_B(t,x_1)-S_{f_0}(t,x_2)],$$

$$E_{ai}^{(q)}(t)=\alpha_s(t)a_i^{(q)}(t)\exp[-S_{f_0}(t,x_2)-S_K(t,x_3)],$$

where the exponential effects denote Sudakov factors generated by the k_T resummation. Since the annihilation topologies in both contents are associated with the matrix elements $\langle f_0(980)K|\bar{s}\gamma_\mu\gamma_5 d(u)|0\rangle$ and $\langle f_0(980)K|\bar{s}\gamma_5 d(u)|0\rangle$, the difference in the different content is only from the spectator, which is s (d or u) in the $\bar{s}s(\bar{n}n)$ component. Hence $N_{a4(6)}^{P(q)}$ could be obtained from $S_{a4(6)}^{P(q)}$ by replacing x_2 with $1-x_2$ in $\Phi_{f_0}(x_2)$ and $1-x_3$ with x_3 in $\{\Phi_K(1-x_3)\}$. In Eq. (11), we define that $S_{a46}^{P(q)}=S_{a4}^{P(q)}+S_{a6}^{P(q)}$ and $N_{e46}^{P(q)}=N_{e4}^{P(q)}+N_{e6}^{P(q)}$. The S_a and N_e can be obtained from $S_{a4}^{P(q)}$ and $N_{e4}^{P(q)}$ by replacing the WC $a_4^{(q)}$ with a_2. We note that unlike the cases of $B\to PP$ and $B\to PV$ decays in which the chirality suppression in the $(V-A)\times(V-A)$ annihilation topology is conspicuous, the contribution of Eq. (13) may not be small because one twist-3 effect, such as Φ_K^σ, is not canceled. Also, differing from $B\to PP$ decays, Eq. (15) is opposite to Eq. (16) in sign.

Similar to $B\to f_0(980)K$ decays, the amplitudes for $B\to f_0(980)K^{*0}$ and $B^+\to f_0(980)K^{*+}$ can be expressed as

$$A_{ss}^{-}=\tilde{f}V_t^*K_{e6}^{P(s)}+f_BV_t^*K_{a46}^{P(d)}+\cdots,$$

$$A_{nn}^{-}=f_{K*}V_t^*X_{e4}^{P(d)}+f_BV_t^*X_{a46}^{P(d)}+\cdots,$$

$$A_{ss}^{\pm}=\tilde{f}V_t^*K_{e6}^{P(s)}+f_BV_t^*K_{a46}^{P(u)}-f_BV_u^*K_a+\cdots,$$

$$A_{nn}^{\pm}=f_{K*}V_t^*X_{e4}^{P(u)}+f_BV_t^*X_{a46}^{P(u)}-f_{K*}V_u^*X_e -f_BV_u^*X_a+\cdots, \quad (17)$$

where $K_{e6}^{P(q)}$, $K_{a4}^{P(q)}$, and $K_{a6}^{P(q)}$ can be obtained from $S_{e6}^{P(q)}$, $S_{a4}^{P(q)}$, and $S_{a6}^{P(q)}$, respectively, by replacing ϕ_K, ϕ_K^p, and ϕ_K^σ with ϕ_{K*}, ϕ_{K*}^s, and ϕ_{K*}^t. With the similar analysis, we find that $X_{e4}=-N_{e4}$ and $X_{a4(6)}^{P(q)}$ are related to $K_{a4(6)}^{P(q)}$ if we change x_2 and $1-x_3$ to $1-x_2$ and x_3, respectively. The K_a and X_e are the same as $K_{a4}^{P(q)}$ and $X_{e4}^{P(q)}$ but the associated WC is a_2.

In our numerical calculations, the B meson wave function is taken as [21]

$$\Phi_B(x,b)=N_Bx^2(1-x)^2\exp\left[-\frac{1}{2}\left(\frac{xM_B}{\omega_B}\right)^2-\frac{\omega_B^2b^2}{2}\right],$$

where ω_B is the shape parameter and N_B is the normalization, determined by

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

Since $f_0(980)$ is a light meson, its wave function can be defined in the frame of the light-cone and the concept of the twist expansion can be used. Because the relevant scalar me-

TABLE I. The parameters in the amplitudes with $\omega_B = 0.4$, $m_K^0 = 1.7$, and $\tilde{f} = 0.2$ GeV for $\xi = 0.3$.

Amplitude	$S_{e6}^{P(s)}(10^{-2})$	$S_{a46}^{P(d)}(10^{-2})$	$S_a(10^{-2})$	$N_{e46}^{P(u)}(10^{-2})$	$N_e(10^{-2})$
$f_0(980)K$	-1.02	$0.45+i0.39$	$-7.41-i0.40$	1.75	31.44
Amplitude	$K_{e6}^{P(s)}(10^{-2})$	$K_{a46}^{P(d)}(10^{-2})$	$K_a(10^{-2})$	$X_{e4}^{P(u)}(10^{-2})$	$X_e(10^{-2})$
$f_0(980)K^*$	-1.37	$0.06+i0.40$	$0.83-i1.98$	1.31	-31.44

son wave function has not been derived in the literature yet, we choose the following form:

$$\Phi_{f_0}(x) = \frac{\tilde{f}}{2\sqrt{2N_c}}\{3(1-2x)^2 + \xi(1-2x)^2[C_2^{3/2}(1-2x)-3] + 1.8C_4^{1/2}(1-2x)\},$$

with $C_4^{1/2}(y) = (35y^4 - 30y^2 + 3)/8$, $C_2^{3/2}(y) = 3/2(5y^2-1)$, and $\xi = 0.3$–0.5 in our estimations. In order to fix the values of ω_B and m_K^0, we take $\omega_B = 0.4$ and $m_K^0 = 1.7$ GeV as those in the studies of the $B\to K$ and $B\to K^*$ form factors [21,27]. Explicitly, with the values above and $\tilde{f} = 0.2$ [7], the $B \to f_0(980)$ form factor is found to be 0.270 (0.286) for $\xi = 0.3$ (0.5). Using Eqs. (11)–(17), the values of hard amplitudes are shown in Table I. In addition, the decay BRs and CP asymmetries ($\mathcal{A}_{CP}$) for $B\to f_0(980)K$ and $B \to f_0(980)K^*$ are shown in Table II with $\phi_s = 0$ and in Table III with $\phi_s = 40°$ and $140°$, where

$$\mathcal{A}_{CP} = \frac{\bar{\Gamma} - \Gamma}{\bar{\Gamma} + \Gamma},$$

with Γ ($\bar{\Gamma}$) being the particle (antiparticle) partial decay rate. In both tables, the CP violating phase of V_{ub} is taken to be $\phi_3 = 72°$. From Table I, we see clearly that the sign of annihilation contribution $S_{a46}^{P(d)}$ is opposite to $S_{e6}^{P(s)}$ so that the BRs of $B\to f_0(980)K$ are reduced significantly. Because the corresponding value of $S_{a46}^{P(d)}$ in $B\to f_0(980)K^*$ is too small as shown in Table II, we can understand that the influence of the annihilation is insignificant. Also we see that if $f_0(980)$ is $\bar{s}s$ mostly, BRs of $B\to f_0(980)K$ are only around 10^{-6}. On the contrary, they could be over 5.0×10^{-6} at a large angle of ϕ_s; whereas BRs for $B\to f_0(980)K^*$ are similar to those at $\phi_s = 0$. To further understand the effect of $\bar{n}n$, we display the BRs as a function of ϕ_s in Fig. 1. Obviously, the predictions are very sensitive to the contributions of the $\bar{n}n$ component. The essential question is what range of ϕ_s is allowed. Unfortunately, the preferable ϕ_s is still unknown, i.e., both small [28] and large [9] angle solutions exist simultaneously in the literature. The former prefers the value of $\phi_s = 42.14^{+5.8^0}_{-7.3}$ but the latter $138^0 \pm 6^0$.

It is also interesting to look for a decay which is directly related to the $\bar{n}n$ content only such that it can tell us how the impact of the content would be. We point out that one of the possible decays is $B^+\to f_0(980)\pi^+$. In this decay, from the identity of Eq. (4) the contribution from the $\bar{s}s$ content vanishes since it corresponds to the vector current vertex. As the π meson is similar to K except $SU(3)$ breaking effects as well as the relevant wave functions and CKM matrix elements, the BR is estimated to be 0.25 $(0.45)\times 10^{-6}$ for $\phi_s = 35°$ $(132°)$.

It seems that in spite of the uncertain allowed value of ϕ_s, the predicted BR of $B^+\to f_0(980)K^+$ is smaller than the central (minimal) value of 21.17 $(11.03)\times 10^{-6}$ reported by Belle in which $\mathrm{Br}[f_0(980)\to \pi^+\pi^-] = 2/3R$ with $R = \Gamma(\pi\pi)/\Gamma(\pi\pi) + \Gamma(KK) \sim 0.68$ [29] is used. It is clear that unless there exists some other mechanism in the decay processes, it is difficult to explain the large measured BR by just considering short-distance interactions. Actually, the case is similar to the charmed decay of $B^+\to \chi_{c0}K^+$ with its BR being $(6.0^{+2.1}_{-1.8}\pm 1.1)\times 10^{-4}$ [30] and $(2.4\pm 0.7)\times 10^{-4}$ [31] measured by Belle and BaBar, respectively. In order to agree with the experimental data, the authors of Ref. [32] suggest that the possible rescattering effects, such as $B \to D_s^* D$, $B\to D_s D^*$, and $B\to D_s^* D^*$ decays via triangle dia-

TABLE II. BRs of $B\to f_0(980)K^{(*)}$ without and with annihilation contributions by taking the same values of parameters as Table I. The *CPA*s with annihilation effects are also shown.

Mode	BR($\times 10^{-6}$) no annihilation	BR($\times 10^{-6}$)	*CPA* (%)
$B\to f_0(980)K^0$	2.95	1.39	0
$B^+\to f_0(980)K^+$	3.13	1.57	6.50
$B\to f_0(980)K^{*0}$	5.16	5.40	0
$B^+\to f_0(980)K^{*+}$	5.49	5.76	1.48

TABLE III. BRs and *CPA*s of $B\to f_0(980)K^{(*)}$ decays with the same parameters as Table I but by taking $\phi_s = 40°$ and $140°$.

Mode	ϕ_s	BR($\times 10^{-6}$)	*CPA* (%)
$B\to f_0(980)K^0$	40°	1.14	0
	140°	4.70	0
$B^+\to f_0(980)K^+$	40°	1.25	-16.94
	140°	5.94	1.58
$B\to f_0(980)K^{*0}$	40°	0.85	0
	140°	6.70	0
$B^+\to f_0(980)K^{*+}$	40°	1.51	11.31
	140°	6.30	-12.19

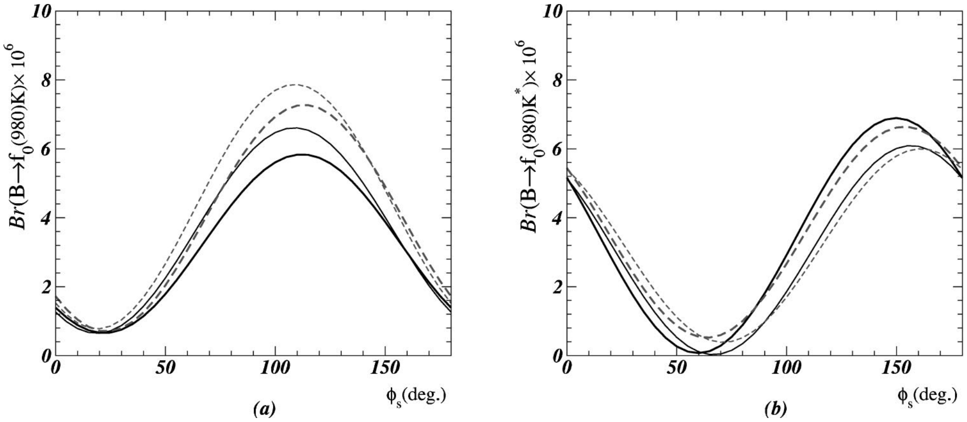

FIG. 1. BRs of (a) $B \to f_0(980)K$ and (b) $B \to f_0(980)K^{(*)}$ as a function of ϕ_s. The bold (thin) solid lines stand for neutral modes with $\xi = 0.3(0.5)$ while the dashed lines are for charged modes.

grams with the strong couplings of $g_{DD\chi_{c0}}$, $g_{D^*D^*\chi_{c0}}$, $g_{D^*D^*\chi_{c0}}$, etc., could enhance the BR and reach (1.1–3.2) $\times 10^{-4}$. We find that the mechanism could also apply to $B \to f_0(980)K^{(*)}$ decays by replacing χ_{c0} with $f_0(980)$ and taking the proper couplings such as $g_{D_s^* D_s f_0(980)}$, etc. The illustrated diagrams are displayed in Fig. 2. Since the differences between $f_0(980)K^{(*)}$ and $\chi_{c0}K^+$ modes are only related to the strong coupling constants, by following the procedure of Ref. [32] and taking proper values of strong coupling constants, one expects that the rescattering effects in Fig. 2 can reach $\mathcal{O}(10^{-5})$ easily. In addition, we note that, as the possibility of a gluonium state in η' [24,33] was proposed to explain the tendency of large BR for η' production [34], it is also possible to understand $B \to f_0(980)K$ decays by considering $g-g-f_0(980)$ coupling.

Finally, it is worth mentioning that like the decays of $B \to J/\psi K_{S,L}$ and $B \to \phi K_{S,L}$, $B \to f_0(980)K_{S,L}$ can be another outstanding candidate for the observation of the time-dependent CP asymmetry, defined by

$$\frac{\overline{\Gamma}(t)-\Gamma(t)}{\overline{\Gamma}(t)+\Gamma(t)} = -A_{CP}^{dir}(B_d \to f)\cos(\Delta M_d t) - A_{CP}^{mix}(B_d \to f)\sin(\Delta M_d), \quad (18)$$

with

$$A_{CP}^{dir}(B_d \to f) = \frac{1-|\lambda|^2}{1+|\lambda|^2},$$

$$A_{CP}^{mix}(B_d \to f) = \frac{2Im\lambda}{1+|\lambda|^2},$$

$$\lambda = \eta e^{-2i\phi_1}\frac{A(\overline{B}_d \to f)}{A(B_d \to f)}, \quad (19)$$

where f expresses the CP final state classified by $\eta = \pm 1$, A_{CP}^{dir} (A_{CP}^{mix}) denotes the direct (mixing-induced) CP violation and $A[B(\overline{B}_d) \to f]$ are the $B(\overline{B})$ meson decay amplitudes and $\phi_1 = \arg(V_{td}^* V_{tb})$. From Eq. (19), we know that if no CP violating phase is involved in the decay amplitude, the direct CP violating observable will vanish and $\lambda = \eta e^{-2i\phi_1}$. Consequently, the mixing-induced CP asymmetry is only related to $\sin 2\phi_1$. Fortunately, since there is no tree contribution, $B \to f_0(980)K_{S,L}$ decays satisfy the criterion so that the mixing-induced CP asymmetry of $B \to f_0(980)K_{S,L}$ is expected to be the same as that measured in $B \to J/\psi K_{S,L}$ decays, except a small derivation from higher order contributions [35].

We have studied $f_0(980)$ scalar meson production in B meson decays by assuming that its flavor contents are $\bar{s}s$ and $\bar{n}n$ states. We have found that the role of $\bar{n}n$ on the BRs of $B \to f_0(980)K$ is crucial. We have also pointed out that non-vanishing BR of $B^+ \to f_0(980)\pi^+$ could test the existence of the $\bar{n}n$ content although the expected BR is less than 10^{-6}. If the BRs of $B \to f_0(980)K^{(*)}$ are measured to be 10^{-5} in experiments, the results could be the evidence of existing rescattering effects in heavy B meson decays. On the other hand, although we concentrate on the $\bar{q}q$ states of $f_0(980)$, we do not exclude the possibilities of four-quark $qq\bar{q}\bar{q}$ and gluonium states. The possibilities could be clarified by other experiments, such as $\phi \to \pi\pi(KK)\gamma$ [36], $D_s^+ \to \pi^+\pi^+\pi^-$ [37] decays, and $B^+ \to f_0(980)\pi^+$ with the measured BR of $\mathcal{O}(10^{-6})$. Furthermore, the more general approach to deal

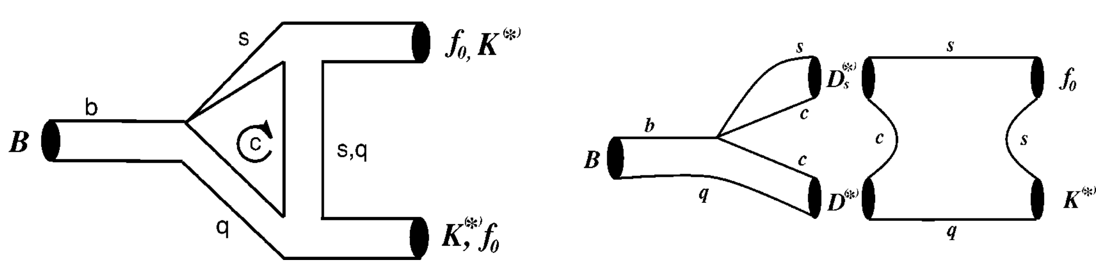

FIG. 2. Possible rescattering diagrams contributing to $B \to f_0(980)K^{(*)}$.

with the decays of $B \to f_0(980)K$ can refer to Ref. [38]. We have also suggested that time-dependent *CP* asymmetry in $B \to f_0(980)K_{S,L}$ could be a new explorer of $\sin 2\phi_1$. We note that our study of B decays with $f_0(980)$ in the final state can be applied to other modes with scalar or isoscalar final states such as $\sigma(600)$, $a_0(980)$, etc. Remarkably, B factories have opened a new opportunity to understand the contents of scalars and their productions.

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