Abstract

We study the exclusive decays of $B \to f_0(980)K(\ast)$ in the framework of the perturbative QCD by identifying the $f_0(980)$ as the composition of $\bar{s}s$ and $\bar{n}n = (\bar{u}u + \bar{d}d)/\sqrt{2}$. We find that the influence of the $\bar{n}n$ content on the predicted branching ratios is crucial. We discuss the possible rescattering and gluonium states which could enhance the branching ratios of considered decays. We point out that the CP asymmetry in $B \to f_0(980)K_{S,L}$ could be a new explorer of $\sin2\phi_1$. 
Recently, Belle \cite{1} and BaBar \cite{2} have opened new channels in three-body nonleptonic $B$ decays, such as $B \rightarrow KKK$, $KK\pi$ and $K\pi\pi$ decays. In particular, BELLE has observed the decay of $B^+ \rightarrow f_0(980)K^+$ with the branching ratio (BR) product of $Br(B^+ \rightarrow f_0(980)K^+) \times Br(f_0(980) \rightarrow \pi^+\pi^-) = (9.6^{+2.5+1.5+3.4}_{-2.3-1.5-0.8}) \times 10^{-6}$ \cite{1}. Since $f_0(980)$ is a neutral scalar meson, the measured $B \rightarrow f_0(980)K$ decays not only show that in the first time $B$ 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. \cite{3} with phase shift analysis. In the literature, $f_0(980)$ could be four-quark states denoted by $qq\bar{q}\bar{q}$ \cite{4} or $KK$ molecular states \cite{5} or $\bar{q}q$ states \cite{6}. However, one objection against the possibility of $KK$ states is that the binding energy of $10-20$ MeV for $KK$ is much smaller than measured width in the range $40-100$ MeV \cite{7}. It is suggested that in terms of the measured $\phi \rightarrow f_0(980)\gamma$ and $f_0(980) \rightarrow \gamma\gamma$ \cite{4} \cite{8} \cite{9} \cite{10} decays and $D_s^+ \rightarrow f_0(980)\pi^+$ decay \cite{11} \cite{12}, the flavor contents of $f_0(980)$ are $\bar{s}s$ mostly and a small portion of $\bar{n}n = (\bar{u}u + \bar{d}d)/\sqrt{2}$. In this paper, we take $f_0(980)$ to be composed of $\bar{q}q$ states and use $|f_0(980) > = \cos \phi_s |\bar{s}s > + \sin \phi_s |\bar{n}n >$ 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 on $B \rightarrow f_0(980)K$. For simplicity, our analysis will only concentrate on dominant factorizable parts and regard the $f_0(980)$ as the composition of $\bar{s}s$ so that at the quark level the process corresponds to $b \rightarrow s\bar{s}s\bar{s}$ decay. Since at the quark level, the physics for $B \rightarrow \phi K$ decays is the same as $B \rightarrow f_0(980)K(\ast)$, we first examine the $\phi K$ mode. We start by writing the effective Hamiltonian for the $b \rightarrow s$ transition as

$$H_{\text{eff}} = \frac{G_F}{\sqrt{2}} \sum_{q' = u, c} V_{q'}^* \left[ C_1(\mu)O_1^{(q')} + C_2(\mu)O_2^{(q')} + \sum_{i=3}^{10} C_i(\mu)O_i \right],$$

where $V_{q'} = V_{q's}^* V_{q'b}$ are the products of the CKM matrix elements, $C_i(\mu)$ are the Wilson coefficients (WCs) and $O_i$ correspond to the four-quark operators. The explicit expressions of $C_i(\mu)$ and $O_i$ can be found in Ref. \cite{12}. It is known that the vector meson $\phi$ has the current matrix elements

$$\langle 0 | \bar{s}\gamma_\mu s | \phi \rangle = m_\phi f_\phi \varepsilon_\mu, \quad \langle 0 | \bar{s}s | \phi \rangle = 0. \quad (1)$$
Therefore, the decay amplitude of $B_d \to \phi K^0$ can be simply described by

$$A \left( B_d \to \phi K^0 \right) = f_\phi V_t^\ast \left( a_3^{(s)} + a_4^{(s)} + a_5^{(s)} \right) F_e^{BK} + 2 f_B V_t^* a_6^{(d)} F_{a6}^{\phi K} + ...$$

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

$$a_1 = C_1 + \frac{C_2}{N_c}, a_2 = C_2 + \frac{C_1}{N_c},$$
$$a_{3,4}^{(q)} = C_{3,4} + \frac{3\epsilon_q}{2} \left( a_{3,4}^{(q)} + a_{3,4}^{(q)} \right), a_{3,4}^{(q)} = \frac{C_{4,3}}{N_c} + \frac{3\epsilon_q}{2N_c} C_{10,9},$$
$$a_{5,6}^{(q)} = C_{5,6} + \frac{3\epsilon_q}{2} C_{7,8} + a_{5,6}^{(q)}, a_{5,6}^{(q)} = \frac{C_{6,5}}{N_c} + \frac{3\epsilon_q}{2N_c} C_{8,7}. \quad (3)$$

Note that $a_2$ is larger than $a_1$ and $a_{3,4}^{(q)}$ 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 $\langle \phi K | \bar{s} \gamma_5 \gamma_d | 0 \rangle$ 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

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

where $m_{f_0}(\tilde{f})$ are the mass (decay constant) of $f_0(980)$. From Eqs. (1) and (2), it is clear that the situation in the $f_0(980)K^0$ mode is just opposite to the case of $\phi K^0$, i.e., the role in vector and scalar vertex is exchanged each other. Hence, the decay amplitude for $B \to f_0(980)K^0$ can be simply written as

$$A \left( B_d \to f_0(980)K^0 \right) = 2r_f \tilde{f} V_t^* a_6^{(d)} S_e^{BK} + 2 f_B V_t^* a_6^{(d)} F_{a6}^{f_0K} + ... \quad (5)$$

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. (3) 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 $a_6^{(s)}(\sqrt{\Lambda_M})/a_4^{(s)}(\sqrt{\Lambda_M}) \approx 1.5$ and $\tilde{f} \sim f_\phi$ and neglecting the annihilation contributions, we estimate the ratio of $Br(B \to f_0(980)K^0)/Br(B \to \phi K^0)$ being around 0.31. With the average value of Belle and BarBar, the predicted $Br(B \to f_0(980)K^0)$ is around $2.8 \times 10^{-6}$. 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 uncertainties are always involved in the calculations of transition matrix elements while studying exclusive hadron decays. Nevertheless, the problem will become mildly 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], \( B \to \phi\pi(K) \) [14, 23], \( B \to \eta^{(')}K \) [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

\[
\langle 0|\bar{q}(0)q(z)|f_0\rangle = \frac{1}{\sqrt{2N_c}} \int_0^1 dx e^{-ixP_z}m_f[1]_{ij}\Phi_{f_0}(x) \tag{6}
\]

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

\[
\langle 0|\bar{b}(z)q(0)|B\rangle = \frac{1}{\sqrt{2N_c}} \int_0^1 dx e^{-ixP_z}\left[[\gamma_5 P]_{jl} + M_B[I]_{jl}\gamma_5\phi_B(x)\right] \tag{7}
\]

\[
\langle K|\bar{q}(z)s(0)|0\rangle = \frac{1}{\sqrt{2N_c}} \int_0^1 dx e^{-ixP_z}\left\{[\gamma_5 P]_{jl}\Phi_K(x) + m_0[K]_{jl}\gamma_5(\not{\epsilon}_- \not{\epsilon}_+ - I)_{jl}\phi_K^s(x)\right\}, \tag{8}
\]

\[
\langle K^*(\epsilon_L)|\bar{q}(z)s(0)|0\rangle = \frac{1}{\sqrt{2N_c}} \int_0^1 dx e^{-ixP_z}\left\{[\gamma_5 P]_{jl}\phi_{K^*}(x) + M_{K^*}[\epsilon_L]_{jl}\phi_{K^*}(x)\right\}, \tag{9}
\]

with \( n_- = (0, 1, 0, \perp) \), \( n_+ = (1, 0, 0, \perp) \) and \( m_0^K \) 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 Ref. [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 \( B \rightarrow f_0(980)K^{(*)} \), we consider only the simplest case. Moreover, if \( f_0(980) \) consists of \( \bar{ss} \) mostly, the choice of Eq. (8) is clearly dominant.

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

\[
\Gamma = \frac{G_F^2 M_B^3}{32\pi} |A|^2
\]  

(10)

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

\[
\begin{align*}
A_{\bar{ss}} &= \tilde{f}_V s_{s_{66}}^{P(s)} + f_B V_s s_{a_{46}}^{P(d)} + \ldots, \\
A_{\bar{nn}} &= f_K V_s s_{a_{46}}^{P(d)} + f_B V_s s_{a_{46}}^{P(d)} + \ldots, \\
A^+_{\bar{ss}} &= \tilde{f}_V s_{s_{66}}^{P(s)} + f_B V_s s_{a_{46}}^{P(u)} - f_B V_s s_{a} + \ldots, \\
A^+_{\bar{nn}} &= f_K V_s s_{a_{46}}^{P(u)} + f_B V_s s_{a_{46}}^{P(u)} - f_K V_s s_{a} - f_B V_s s_{a} + \ldots, 
\end{align*}
\]  

(11)

respectively, where \( s_{e_{(a)}}^{P(q)} (n_{e_{(a)}}^{P(q)}) \) denote the emission (annihilation) contributions of the \( \bar{ss} \) (\( \bar{nn} \)) 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 \( A = \cos \phi A_{\bar{ss}}^{(+)} + \sin \phi A_{\bar{nn}}^{(+)} / \sqrt{2} \). Due to the smallness of nonfactorizable effects, we neglect to show them in the 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. (13-15), the hard amplitudes for \( \bar{ss} \) content are derived as

\[
\begin{align*}
S_{e_{66}}^{P(q)} &= 16\pi r_f \alpha_s C_F M_B^2 \int_0^1 dx_1 dx_3 \int_0^\infty b_1 db_1 b_2 db_3 \Phi_B(x_1, b_1) \\
&\times \left\{ \left[ \Phi_K(x_2) + 2 r_K \Phi_{K}^P(x_3) + r_K x_3 (\Phi_{K}^P(x_3) - \Phi_{K}^P(x_3)) \right] \\
&\times E S_{e_6}^{(q)} (t_e^{(1)}) h_e(x_1, x_3, b_1, b_3) \\
&+ 2 r_K \Phi_{K}^P(x_3) E S_{e_6}^{(q)} (t_e^{(2)}) h_e(x_1, x_3, b_1, b_3) \right\}, \\
S_{a_{44}}^{P(q)} &= -16\pi r_f r_K C_F M_B^2 \int_0^1 dx_2 dx_3 \int_0^\infty b_2 db_2 b_3 db_3 \Phi_{f_0}(x_2) \\
&\times \left\{ \left[ (1 + x_3) \Phi_{K}^P(1 - x_3) + (1 - x_3) \Phi_{K}^P(1 - x_3) \right] E_{a_{44}}^{(q)} (t_a^{(1)}) h_a(x_3, x_2, b_2, b_3) \\
&- \left[ (1 + x_2) \Phi_{K}^P(1 - x_3) \right] E_{a_{44}}^{(q)} (t_a^{(2)}) h_a(x_3, x_2, b_3, b_2) \right\},
\end{align*}
\]  

(12-13)

5
\[ S_{a6}^{P(q)} = 16\pi r_f C_F M_B^2 \int_0^1 dx_2 dx_3 \int_0^\infty b_2 db_2 b_3 \Phi_f(x_2) \Phi_K(1 - x_3) \times \left\{ 2E_{a6}^{(q)}(t_a^{(1)}) h_a(x_2, x_3, b_2, b_3) + x_2 E_{a6}^{(q)}(t_a^{(2)}) h_a(x_3, x_2, b_3, b_2) \right\} \]  

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

\[ N_{e4}^{P(q)} = 8\pi r_f \alpha_s C_F M_B^2 \int_0^1 dx_1 dx_2 \int_0^\infty b_1 db_1 b_2 db_2 \Phi_B(x_1, b_1) \Phi_f(x_2) \times \left\{ (1 - 2x_2) E_{e4}^{(q)}(t_e^{(1)}) h_e(x_1, x_2, b_1, b_2) + 2EN_{e6}^{(q)}(t_e^{(2)}) h_e(x_2, x_1, b_2, b_1) \right\} \]  

\[ N_{e6}^{P(q)} = -16\pi r_f \alpha_s C_F M_B^2 \int_0^1 dx_1 dx_2 \int_0^\infty b_1 db_1 b_2 db_2 \Phi_B(x_1, b_1) \Phi_f(x_2), \]

\[ \times \left\{ (2 + x_2) E_{e6}^{(q)}(t_e^{(1)}) h_e(x_1, x_2, b_1, b_2) + 2EN_{e6}^{(q)}(t_e^{(2)}) h_e(x_2, x_1, b_2, b_1) \right\}. \]  

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

\[ E S_{e6}^{(q)}(t) = \alpha_s(t) a_{6}^{(q)}(t) \exp[-S_B(t, x_1) - S_K(t, x_3)], \]

\[ EN_{e6}^{(q)}(t) = \alpha_s(t) a_{4}^{(q)}(t) \exp[-S_B(t, x_1) - S_f(t, x_2)], \]

\[ E_{a6}^{(q)}(t) = \alpha_s(t) a_{4}^{(q)}(t) \exp[-S_f(t, x_2) - S_K(t, x_3)]. \]

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 \( < f_0(980)K|\bar{s}\gamma_\mu\gamma_5d(u)|0 > \) and \( < f_0(980)K|\bar{s}\gamma_\mu d(u)|0 > \), the difference in the different content is only from the spectator, which is \( s \) (d or u) in \( \bar{s}\bar{s} \) (\( \bar{n}n \)) 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 \( \Phi_f(x_2) \) and \( 1 - x_3 \) with \( x_3 \) in \( \{ \Phi_K(1 - x_3) \} \). In Eq. (13), we define that \( S_{a6}^{P(q)} = S_{a4}^{P(q)} + S_{a6}^{P(q)} \) and \( N_{e6}^{P(q)} = N_{e4}^{P(q)} + N_{e6}^{P(q)} \). The \( S_a \) and \( N_e \) can be obtained from \( S_{a4}^{P(q)} \) and \( N_{e4}^{P(q)} \) by replacing W.C. \( a_4^{(q)} \) 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 \( \Phi_K^c \), is not cancelled. And also, differing from \( B \to PP \) decays, Eq. (13) 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

\[ A_{\bar{s}s} = \bar{f}_Vt^*K_s^{P(s)} + f_{B^+}V_{K}^{P(d)} + \ldots, \]

\[ A_{\bar{n}n} = f_KV_tX_{K}^{P(d)} + f_{B^+}V_{X}^{P(d)} + \ldots, \]
\[ A_{ss}^+ = \bar{f}V_t^* K^{P(s)}_{e6} + f_B V_t^* K^{P(u)}_{a46} - f_B V_u^* K_a + \ldots, \]
\[ A_{tm}^+ = f_K V_t^* X^{P(u)}_{e4} + f_B V_t^* X^{P(u)}_{a46} - f_K V_u^* X_e - f_B V_u^* X_a + \ldots \]

where \( K^{P(q)}_{e6} \), \( K^{P(q)}_{a4} \) and \( K^{P(q)}_{a6} \) can be obtained from \( S^{P(q)}_{e6} \), \( S^{P(q)}_{a4} \) and \( S^{P(q)}_{a6} \), respectively, by replacing \( \phi_K \), \( \phi_K^0 \) and \( \phi_K^* \) with \( \phi_K^{*0} \), \( \phi_K^{*1} \) and \( \phi_K^{*2} \). With the similar analysis, we find that \( X_{e4} = -N_{e4} \) and \( X^{P(q)}_{a4(6)} \) are related to \( K^{P(q)}_{a4(6)} \) 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 \( K^{P(q)}_{a4} \) and \( X^{P(q)}_{e4} \) but the associated WC is \( a_2 \).

In our numerical calculations, the \( B \) meson wave function is taken as \[ \Phi_B(x, b) = N_B x^2 (1 - x)^2 \exp \left[ -\frac{1}{2} \left( \frac{x M_B}{\omega_B} \right)^2 - \frac{\omega_B^2 b^2}{2} \right] \]
where \( \omega_B \) 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. And because the relevant scalar meson wave function has not been derived in the literature yet, we choose the following form
\[ \Phi_{f_0}(x) = \frac{\tilde{f}}{2\sqrt{2N_c}} \left\{ 3(1 - 2x)^2 + \xi(1 - 2x)^2(C_2^{3/2}(1 - 2x) - 3) + 1.8C_4^{1/2}(1 - 2x) \right\} \]
with \( C_4^{1/2}(y) = (35y^4 - 30y^2 + 3)/8 \), \( C_2^{3/2}(y) = 3/2(5y^2 - 1) \) and \( \xi = 0.3 \sim 0.5 \) in our estimations. In order to fix the values of \( \omega_B \) and \( m_{f_0}^0 \), we take \( \omega_B = 0.4 \) and \( m_{f_0}^0 = 1.7 \) GeV as those in the studies of the \( B \rightarrow K \) and \( B \rightarrow K^* \) form factors [21, 27]. Explicitly, with the values above and \( \tilde{f} = 0.2 \) [7], the \( B \rightarrow f_0(980) \) 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 (\( A_{CP} \)) for \( B \rightarrow f_0(980)K \) and \( B \rightarrow f_0(980)K^* \) are shown in Table I with \( \phi_s = 0 \) and in Table II with \( \phi_s = 40^0 \) and \( 140^0 \), where
\[ A_{CP} = \frac{\Gamma - \Gamma}{\Gamma + \Gamma} \]
with \( \Gamma (\tilde{\Gamma}) \) being the particle (antiparticle) partial decay rate. In both tables, the CP violating phase of \( V_{ub} \) is taken to be \( \phi_3 = 72^0 \). From Table I, we see clearly that the sign of annihilation contribution \( S^{P(d)}_{a46} \) is opposite to \( S^{P(s)}_{e6} \) so that the BRs of \( B \rightarrow f_0(980)K \) are reduced significantly. Because the corresponding value of \( S^{P(d)}_{a46} \) in \( B \rightarrow f_0(980)K^* \) is
TABLE I: The parameters in the amplitudes with $\omega_B = 0.4$, $m_K^0 = 1.7$ and $\bar{f} = 0.2$ GeV for $\xi = 0.3$.

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

too small as shown in Table I, 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 $B \rightarrow f_0(980)K$ are only around $10^{-6}$. On the contrary, they could be over $5.0 \times 10^{-6}$ at a large angle of $\phi_s$; whereas BRs for $B \rightarrow f_0(980)K^*$ 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 $\phi_s$ in Figure I. Obviously, the predictions are very sensitive to the contributions of the $\bar{n}n$ component. The essential question is what the range of $\phi_s$ is allowed. Unfortunately, the preferable $\phi_s$ is still unknown, i.e., both small [28] and large [9] angle solutions exist simultaneously in the literature. The former prefers the value of $\phi_s = 42.14 \pm 5.8^0$ 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 $B^+ \rightarrow f_0(980)\pi^+$. 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 $\pi$ 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 $\phi_s = 35^0 (132^0)$.

It seems that in spite of the uncertain allowed value of $\phi_s$, the predicted BR of $B^+ \rightarrow f_0(980)K^+$ is smaller than the central (minimal) value of $21.17 (11.03) \times 10^{-6}$ reported by Belle in which $Br(f_0(980) \rightarrow \pi^+\pi^-) = 2/3R$ 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-distant interactions. Actually, the case is similar to the charmed decay of $B^+ \rightarrow \chi_{c0}K^+$ 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.
FIG. 1: BRs of (a) $B \to f_0(980)K$ and (b) $B \to f_0(980)K^{(*)}$ as a function of $\phi_s$. The bold (thin) solid lines stand for neutral modes with $\xi = 0.3(0.5)$ while the dashed lines are for charged modes.

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 CPAs with annihilation effects are also shown.

| Mode               | BR($\times 10^{-6}$) | CPA (%) |
|--------------------|----------------------|---------|
| $B \to f_0(980)K^0$| 2.95                 | 1.39    |
| $B^+ \to f_0(980)K^+$| 3.13                | 1.57    |
| $B \to f_0(980)K^{*0}$| 5.16                | 5.40    |
| $B^+ \to f_0(980)K^{*+}$| 5.49                | 5.76    |

In order to agree with the experimental data, the authors of Ref. [32] suggest that the possible rescattering effects, such as $B \to D_s^*D$, $B \to D_sD^*$ and $B \to D_s^*D^*$ decays via triangle diagrams with the strong couplings of $g_{D_sD_s\chi_c^0}$, $g_{D^*D^*\chi_c^0}$, $g_{D^*D^*\chi_c^0}$ etc., could enhance the BR and reach $(1.1 \sim 3.2) \times 10^{-4}$. We find that the mechanism could also apply to $B \to f_0(980)K^{(*)}$ decays by replacing $\chi_c^0$ with $f_0(980)$ and taking the proper couplings such as $g_{D^*_sD_sf_0(980)}$ etc.. The illustrated diagrams are displayed in Figure 2. Since the differences between $f_0(980)K^{(*)}$ and $\chi_c^0K^+$ 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 Figure 2 can reach $O(10^{-5})$ easily. In addition, we note that, as the possibility of gluonium state in $\eta'$ [24, 33] was proposed to explain the
TABLE III: BRs and CPAs of $B \to f_0(980)K^{(*)}$ decays with the same parameters as Table II but by taking $\phi_s = 40^0$ and $140^0$.

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

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

tendency of large BR for $\eta'$ 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 to mention 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

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

with

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

$$A_{CP}^{mix}(B_d \to f) = \frac{2Im\lambda}{1 + |\lambda|^2}$$
\[ \lambda = \eta e^{-2i\phi_1} \frac{A(\bar{B}_d \to f)}{A(B_d \to f)}, \]  

(19)

where \( f \) expresses the CP final state classified by \( \eta = \pm 1 \), \( A_{\text{dir}}^{CP} \) \((A_{\text{mix}}^{CP})\) denotes the direct (mixing-induced) CP violation and \( A(B(\bar{B}_d) \to f) \) are the \( B(\bar{B}) \) meson decay amplitudes and \( \phi_1 = \text{Arg}(V_{td}^*V_{tb}) \). 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 \( \bar{ss} \) and \( \bar{nn} \) states. We have found the role of \( \bar{nn} \) on the BRs of \( B \to f_0(980)K \) is crucial. We have also pointed out that nonvanishing BR of \( B^+ \to f_0(980)\pi^+ \) could test the existence of the \( \bar{nn} \) 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 \( \bar{qq} \) 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], \( D_\ast^+ \to \pi^+\pi^+\pi^- \) [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 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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