Charmless B decays to a scalar meson and a vector meson

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The hadronic charmless B decays into a scalar meson and a vector meson are studied within the framework of QCD factorization. The main results are: (i) The decay rates for the $f_0(980)K^{*-}$ and $f_0(980)\bar{K}^{*0}$ modes depend on the $f_0 - \sigma$ mixing angle θ. (ii) If the $a_0(980)$ is a $q\bar{q}$ bound state, the predicted branching ratios for the color-allowed channels $a_0^+\rho^-$ and $a_0^0\rho^-$ will be very large, of order 28×10^{-6} and 21×10^{-6}, respectively. (iii) For the $a_0(1450)$ channels, the color-allowed modes $a_0^+(1450)\rho^-$ and $a_0^0(1450)\rho^-$ are found to have branching ratios of order 20×10^{-6} and 30×10^{-6}, respectively. A measurement of them at the predicted level will favor the $q\bar{q}$ structure for the $a_0(1450)$. (iv) Contrary to the naive expectation that $\Gamma(B^- \to a_0^0\rho^-) \sim \frac{1}{2}\Gamma(\bar{B}^0 \to a_0^+\rho^-)$, we found that this naive relation is violated especially for $a_0 = a_0(1450)$ due to additional contributions to the $a_0^0\rho^-$ mode from the a_0^0 emission. (v) The decays $B \to K_0^*(1430)\rho$ are expected to have rates substantially larger than that of $B \to K_0^*(1430)\pi$ owing to the constructive (destructive) interference between the a_4 and a_6 penguin terms in the former (latter). Experimentally, it is thus important to check if the $B \to K_0^*\rho$ modes are enhanced relative to the corresponding $K_0^*\pi$ channels.

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I. INTRODUCTION

Recently we have studied the hadronic charmless B decays into a scalar meson and a pseudoscalar meson within the framework of QCD factorization (QCDF) [1]. It is known that the identification of scalar mesons is difficult experimentally and the underlying structure of scalar mesons is not well established theoretically (for a review, see e.g. [2–4]). The experimental measurements of $B \to SP$ will provide valuable information on the nature of the even-parity mesons. For example, it was claimed by us [1] that the predicted $\bar{B}^0 \to a_0^\pm(980)\pi^\mp$ and $a_0^+(980)K^-$ rates exceed the current experimental limits, favoring a four-quark nature for the $a_0(980)$. The decay $\bar{B}^0 \to \kappa^+K^-$ also provides a nice ground for testing the 4-quark and 2-quark structure of the κ [or $K_0^*(800)$] meson. It can proceed through W-exchange and hence is quite suppressed if the κ is made of $q\bar{q}$ quarks, while it receives a tree contribution if the κ is predominately a four-quark state. Hence, an observation of this channel at the level of $\sim 10^{-7}$ may imply a four-quark assignment for the κ [1].

In this work we shall generalize our previous study to the decays $B \to SV$ (S: scalar meson, V: vector meson), motivated by the recent observation of the $\bar{K}_0^{*0}(1430)\phi$ and $f_0(980)K^{*-}$ modes by *BABAR* [5–9]:

$$
\begin{aligned}
\mathcal{B}(\bar{B}^0 \to \bar{K}_0^{*0}(1430)\phi) &= (4.6 \pm 0.7 \pm 0.6) \times 10^{-6},\\
\mathcal{B}(B^- \to f_0(980)K^{*-}; f_0(980) \to \pi^+\pi^-) &= (5.2 \pm 1.2 \pm 0.5) \times 10^{-6},\\
\mathcal{B}(\bar{B}^0 \to f_0(980)\bar{K}^{*0}; f_0(980) \to \pi^+\pi^-) &= (2.6 \pm 0.6 \pm 0.9) \times 10^{-6} < 4.3 \times 10^{-6},\\
\mathcal{B}(B^- \to f_0(980)\rho^-; f_0(980) \to \pi^+\pi^-) &< 1.9 \times 10^{-6},\\
\mathcal{B}(\bar{B}^0 \to f_0(980)\rho^0; f_0(980) \to \pi^+\pi^-) &< 0.53 \times 10^{-6},\\
\mathcal{B}(\bar{B}^0 \to f_0(980)\omega; f_0(980) \to \pi^+\pi^-) &< 1.5 \times 10^{-6}.
\end{aligned}
\tag{1.1}
$$

Recently, the decay $\bar{B}^0 \to \bar{K}_0^{*0}(1430)\phi$ has been studied in [10] within the framework of generalized factorization in which the nonfactorizable effects are described by the parameter $N_c^{\rm eff}$, the effective number of colors. The result is sensitive to $N_c^{\rm eff}$. For example, the branching ratio is predicted to be $(7.70, 3.95, 1.84) \times 10^{-6}$ for $N_c^{\rm eff} = (2, 3, 5)$. Hence, in the absence of information for nonfactorizable effects, one cannot make a precise prediction of its rate. A QCDF calculation of this and other modes will be presented in this work.

Since $B \to SP$ decays have been systematically explored in [1], it is straightforward to generalize the study to the SV modes. In the sector of odd-parity mesons, it is known that the rates of the penguin-dominated modes $K^*\pi$ and $K\rho$ are smaller than that of the corresponding $K\pi$ ones by a factor of ~ 2. This can be understood as follows. In the factorization approach, the penguin terms a_6 and a_8 are absent in the decay amplitudes of $B \to K^*\pi$, while the effective Wilson coefficients a_4 and a_6 contribute destructively to $B \to K\rho$. In contrast, the tree-dominated $\rho\pi$

modes have rates larger than that of $\pi\pi$ with the same charge assignment due mainly to the fact that the ρ meson has a decay constant larger than the pion. We shall see in the present work that the same analog is not always true in the scalar meson sector. For example, we will show that the rates for $\bar{K}_0^{*0}(1430)\rho^{-,0}$ are larger than that of $\bar{K}_0^{*0}(1430)\pi^{-,0}$.

The layout of the present paper is as follows. In Sec. II we introduce the input quantities relevant to the present work, such as the decay constants, form factors, and light-cone distribution amplitudes. We then apply QCD factorization in Sec. III to study $B \to SV$ decays. Results and discussions are presented in Sec. IV. Section V contains our conclusions. The factorizable amplitudes of various $B \to SV$ decays are summarized in Appendix A.

II. INPUT QUANTITIES

Since most of the essential input quantities are already discussed in [1], here we shall just recapitulate the main inputs.

A. Decay constants and form factors

Decay constants of scalar and vector mesons are defined as

$$\begin{aligned} \langle V(p)|\bar{q}_2\gamma_\mu q_1|0\rangle &= f_V m_V \varepsilon^*_\mu, \\ \langle S(p)|\bar{q}_2\gamma_\mu q_1|0\rangle &= f_S p_\mu, \\ \langle S|\bar{q}_2 q_1|0\rangle &= m_S \bar{f}_S. \end{aligned} \tag{2.1}$$

For vector mesons, there is an additional transverse decay constant defined by

$$\langle V(p,\varepsilon^*)|\bar{q}\sigma_{\mu\nu}q'|0\rangle = f_V^\perp(p_\mu\varepsilon^*_\nu - p_\nu\varepsilon^*_\mu), \tag{2.2}$$

which is scale dependent. The neutral scalar mesons σ, f_0, and a_0^0 cannot be produced via the vector current owing to charge conjugation invariance or conservation of vector current:

$$f_\sigma = f_{f_0} = f_{a_0^0} = 0. \tag{2.3}$$

For other scalar mesons, the vector decay constant f_S and the scale-dependent scalar decay constant $\bar{f}_S$ are related by equations of motion,

$$\mu_S f_S = \bar{f}_S, \qquad \text{with } \mu_S = \frac{m_S}{m_2(\mu) - m_1(\mu)}, \tag{2.4}$$

where m_2 and m_1 are the running current quark masses and m_S is the scalar meson mass. For the neutral scalar mesons f_0, a_0^0 and σ, f_S vanishes, but the quantity $\bar{f}_S = f_S\mu_S$ remains finite.

In [1] we have considered two different scenarios for the scalar mesons above 1 GeV, which will be briefly discussed in Sec. IV. In the same work we have applied the QCD sum rule method to estimate various decay constants for scalar mesons which are summarized as follows:

$$\begin{aligned} \bar{f}_{a_0(980)}(1\text{ GeV}) &= (365 \pm 20)\text{ MeV}, \\ \bar{f}_{a_0(980)}(2.1\text{ GeV}) &= (450 \pm 25)\text{ MeV}, \\ \bar{f}_{f_0(980)}(1\text{ GeV}) &= (370 \pm 20)\text{ MeV}, \\ \bar{f}_{f_0(980)}(2.1\text{ GeV}) &= (460 \pm 25)\text{ MeV}, \\ \bar{f}_{a_0(1450)}(1\text{ GeV}) &= (460 \pm 50)\text{ MeV}, \\ \bar{f}_{a_0(1450)}(2.1\text{ GeV}) &= (570 \pm 60)\text{ MeV}, \\ \bar{f}_{K_0^*(1430)}(1\text{ GeV}) &= -(300 \pm 30)\text{ MeV}, \\ \bar{f}_{K_0^*(1430)}(2.1\text{ GeV}) &= -(370 \pm 35)\text{ MeV}, \end{aligned} \tag{2.5}$$

in scenario 1 and

$$\begin{aligned} \bar{f}_{K_0^*(1430)}(1\text{ GeV}) &= (445 \pm 50)\text{ MeV}, \\ \bar{f}_{K_0^*(1430)}(2.1\text{ GeV}) &= (550 \pm 60)\text{ MeV}, \end{aligned} \tag{2.6}$$

in scenario 2. Using the running quark masses given in Eq. (A13), we obtain the scale-independent decay constants, for example, $f_{a_0(980)^\pm} = 1.0$ MeV in scenario 1 and $f_{a_0(1450)^\pm} = 5.3$ MeV, $f_{K_0^*(1430)} = 35.9$ MeV in scenario 2. For longitudinal and transverse decay constants of the vector mesons, we use (in units of MeV)

$$\begin{aligned} &f_\rho = 216 \pm 3, \qquad f_\omega = 187 \pm 5, \qquad f_{K^*} = 220 \pm 5, \\ &f_\phi = 215 \pm 5, \qquad f_\rho^\perp = 165 \pm 9, \qquad f_\omega^\perp = 151 \pm 9, \\ &f_{K^*}^\perp = 185 \pm 10, \qquad f_\phi^\perp = 186 \pm 9, \end{aligned} \tag{2.7}$$

where the values of f_V and $f_V^\perp$ are taken from [11].

Form factors for $B \to S, V$ transitions are defined by [12]

$$\begin{aligned} \langle V(p')|V_\mu|B(p)\rangle &= -\frac{1}{m_B + m_V}\epsilon_{\mu\nu\alpha\beta}\varepsilon^{*\nu}P^\alpha q^\beta V^{PV}(q^2), \\ \langle V(p')|A_\mu|B(p)\rangle &= i\Big\{(m_B + m_V)\varepsilon^*_\mu A_1^{PV}(q^2) \\ &\quad - \frac{\varepsilon^*\cdot P}{m_B + m_V}P_\mu A_2^{PV}(q^2) \\ &\quad - 2m_V\frac{\varepsilon^*\cdot P}{q^2}q_\mu[A_3^{PV}(q^2) - A_0^{PV}(q^2)]\Big\}, \\ \langle S(p')|A_\mu|B(p)\rangle &= -i\Big[\Big(P_\mu - \frac{m_B^2 - m_S^2}{q^2}q_\mu\Big)F_1^{BS}(q^2) \\ &\quad + \frac{m_B^2 - m_S^2}{q^2}q_\mu F_0^{BS}(q^2)\Big], \end{aligned} \tag{2.8}$$

where $P_\mu = (p + p')_\mu$, $q_\mu = (p - p')_\mu$. As shown in [13], a factor of $(-i)$ is needed in $B \to S$ transition in order to obtain positive $B \to S$ form factors. This also can be checked from heavy quark symmetry [13].

Various form factors for $B \to S, V$ transitions have been evaluated in the relativistic covariant light-front quark model [13]. In this model form factors are first calculated in the spacelike region and their momentum dependence is

TABLE I. Form factors for $B \to \rho$, K^*, $a_0(1450)$, $K_0^*(1430)$ transitions obtained from the covariant light-front model [13].

F	$F(0)$	$F(q^2_{\rm max})$	a	b	F	$F(0)$	$F(q^2_{\rm max})$	a	b
$F_1^{Ba_0(1450)}$	0.26	0.68	1.57	0.70	$F_0^{Ba_0(1450)}$	0.26	0.35	0.55	0.03
$F_1^{BK_0^*(1430)}$	0.26	0.70	1.52	0.64	$F_0^{BK_0^*(1430)}$	0.26	0.33	0.44	0.05
$V^{B\rho}$	0.27	0.79	1.84	1.28	$A_0^{B\rho}$	0.28	0.76	1.73	1.20
V^{BK^*}	0.31	0.96	1.79	1.18	$A_0^{BK^*}$	0.31	0.87	1.68	1.08

fitted to a 3-parameter form:

$$F(q^2) = \frac{F(0)}{1 - a(q^2/m_B^2) + b(q^2/m_B^2)^2}. \tag{2.9}$$

The parameters a, b, and $F(0)$ are first determined in the spacelike region. This parametrization is then analytically continued to the timelike region to determine the physical form factors at $q^2 \geq 0$. The results relevant for our purposes are summarized in Table I. The form factors for B to $f_0(980)$ and $a_0(980)$ transitions are taken to be 0.25 at $q^2 = 0$ [1].

We need to pay a special attention to the decay constants and form factors for the $f_0(980)$. The quark structure of the light scalar mesons below or near 1 GeV has been quite controversial, though it is commonly believed that the $f_0(980)$ and the $a_0(980)$ are primarily four-quark or molecular states. In this work we shall consider the conventional $q\bar{q}$ assignment for the light scalars $f_0(980)$ and $a_0(980)$ as we cannot make predictions for four-quark bound states in QCDF. In the naive quark model, the flavor wave functions of the $f_0(980)$ and $\sigma(600)$ read

$$\sigma = \frac{1}{\sqrt{2}}(u\bar{u} + d\bar{d}), \qquad f_0 = s\bar{s}, \tag{2.10}$$

where the ideal mixing for f_0 and σ has been assumed. In this picture, $f_0(980)$ is purely an $s\bar{s}$ state. However, there also exist some experimental evidences indicating that $f_0(980)$ is not purely an $s\bar{s}$ state (see [14] for details). Therefore, isoscalars $\sigma(600)$ and f_0 must have a mixing

$$\begin{aligned} |f_0(980)\rangle &= |s\bar{s}\rangle\cos\theta + |n\bar{n}\rangle\sin\theta, \\ |\sigma(600)\rangle &= -|s\bar{s}\rangle\sin\theta + |n\bar{n}\rangle\cos\theta, \end{aligned} \tag{2.11}$$

with $n\bar{n} \equiv (\bar{u}u + \bar{d}d)/\sqrt{2}$. Experimental implications for the $f_0 - \sigma$ mixing angle have been discussed in detail in [14]. In this work, we shall use $\theta = 15°$, which is favored by the phenomenological analysis of $B \to f_0 K^*$ decays (see below). In the decay amplitudes involving the $f_0(980)$, we will use the superscripts $q = u, d, s$, to indicate that it is the q quark content of the $f_0(980)$ that gets involved in the interaction. For example, $\bar{f}^s_{f_0} = \bar{f}_{f_0}\cos\theta$ and $F_1^{Bf_0^u} = F_1^{Bf_0}\sin\theta/\sqrt{2}$.

B. Distribution amplitudes

The twist-2 light-cone distribution amplitude (LCDA) $\Phi_S(x)$ and twist-3 $\Phi_S^s(x)$ and $\Phi_S^\sigma(x)$ for the scalar meson S respect the normalization conditions,

$$\begin{aligned} &\int_0^1 dx\Phi_S(x) = f_S, \\ &\int_0^1 dx\Phi_S^s(x) = \int_0^1 dx\Phi_S^\sigma(x) = \bar{f}_S. \end{aligned} \tag{2.12}$$

In general, the twist-2 light-cone distribution amplitude Φ_S has the form

$$\Phi_S(x,\mu) = f_S 6x(1-x)\left[1 + \mu_S \sum_{m=1}^{\infty} B_m(\mu) C_m^{3/2}(2x-1)\right], \tag{2.13}$$

where B_m are Gegenbauer moments and $C_m^{3/2}$ are the Gegenbauer polynomials. For the neutral scalar mesons f_0, a_0^0 σ, only odd Gegenbauer polynomials contribute. In [1] we have applied the QCD sum rules to determine the Gegenbauer moments B_1 and B_3 (see Table II). For twist-3 LCDAs, we use

$$\Phi_S^s(x) = \bar{f}_S, \qquad \Phi_S^\sigma(x) = \bar{f}_S 6x(1-x). \tag{2.14}$$

TABLE II. Gegenbauer moments B_1 and B_3 in scenario 1 (top) and scenario 2 (bottom) at the scales $\mu = 1$ and 2.1 GeV (shown in parentheses) obtained using the QCD sum rule method [1].

State	B_1	B_3
$a_0(980)$	$-0.93 \pm 0.10(-0.64 \pm 0.07)$	$0.14 \pm 0.08(0.08 \pm 0.04)$
$f_0(980)$	$-0.78 \pm 0.08(-0.54 \pm 0.06)$	$0.02 \pm 0.07(0.01 \pm 0.04)$
$a_0(1450)$	$0.89 \pm 0.20(0.62 \pm 0.14)$	$-1.38 \pm 0.18(-0.81 \pm 0.11)$
$K_0^*(1430)$	$0.58 \pm 0.07(0.39 \pm 0.05)$	$-1.20 \pm 0.08(-0.70 \pm 0.05)$
$a_0(1450)$	$-0.58 \pm 0.12(-0.40 \pm 0.08)$	$-0.49 \pm 0.15(-0.29 \pm 0.09)$
$K_0^*(1430)$	$-0.57 \pm 0.13(-0.39 \pm 0.09)$	$-0.42 \pm 0.22(-0.25 \pm 0.13)$

For vector mesons, the normalization for the twist-2 function Φ_V and the twist-3 function Φ_v is given by [15]

$$\int_0^1 dx \Phi_V(x) = f_V, \qquad \int_0^1 dx \Phi_v(x) = 0, \tag{2.15}$$

where the definition for $\Phi_v(x)$ can be found in [15]. The general expressions of these LCDAs read

$$\Phi_V(x,\mu) = 6x(1-x) f_V \left[1 + \sum_{n=1}^{\infty} \alpha_n^V(\mu) C_n^{3/2}(2x-1) \right], \tag{2.16}$$

and

$$\Phi_v(x,\mu) = 3 f_V^{\perp} \left[2x - 1 + \sum_{n=1}^{\infty} \alpha_{n,\perp}^V(\mu) P_{n+1}(2x-1) \right], \tag{2.17}$$

where $P_n(x)$ are the Legendre polynomials. The Gegenbauer moments α_n^V and $\alpha_{n,\perp}^V$ have been studied using the QCD sum rule method. Here we employ the most recent updated values evaluated at $\mu = 1$ GeV [16]:

$$\begin{aligned} \alpha_1^{K^*} &= 0.03 \pm 0.02, & \alpha_{1,\perp}^{K^*} &= 0.04 \pm 0.03, \\ \alpha_2^{K^*} &= 0.11 \pm 0.09, & \alpha_{2,\perp}^{K^*} &= 0.10 \pm 0.08, \\ \alpha_2^{\rho,\omega} &= 0.15 \pm 0.07, & \alpha_{2,\perp}^{\rho,\omega} &= 0.14 \pm 0.06, \\ \alpha_2^{\phi} &= 0.18 \pm 0.08, & \alpha_{2,\perp}^{\phi} &= 0.14 \pm 0.07, \end{aligned} \tag{2.18}$$

and $\alpha_1^V = 0$, $\alpha_{1,\perp}^V = 0$ for $V = \rho, \omega, \phi$.

As stressed in [1], it is most suitable to define the LCDAs of scalar mesons including decay constants. However, in order to make connections between $B \to SV$ and $B \to VV$ amplitudes, it is more convenient to factor out the decay constants in the LCDAs and put them back in the appropriate places. In the ensuing discussions, we will use the LCDAs with the decay constants f_S, $\bar{f}_S$, f_V, $f_V^{\perp}$, f_P being factored out.

III. $B \to SV$ DECAYS IN QCD FACTORIZATION

We shall use the QCD factorization approach [15,17] to study the short-distance contributions to the $B \to SV$ decays with $S = f_0(980)$, $a_0(980)$, $a_0(1450)$, $K_0^*(1430)$, and $V = \rho$, K^* ϕ, ω In QCD factorization, the factorizable amplitudes of above-mentioned decays are summarized in Appendix A. The effective parameters a_i^p with $p = u, c$ in Eq. (A8) can be calculated in the QCD factorization approach [17]. 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 [15,17][1]

$$\begin{aligned} a_i^p(M_1 M_2) = &\left(c_i + \frac{c_{i\pm1}}{N_c} \right) N_i(M_2) \int_0^1 \Phi_{\parallel}^{M_2}(x) dx \\ &+ \frac{c_{i\pm1}}{N_c} \frac{C_F \alpha_s}{4\pi} \left[V_i(M_2) + \frac{4\pi^2}{N_c} H_i(M_1 M_2) \right] \\ &+ P_i^p(M_2), \end{aligned} \tag{3.1}$$

where $i = 1, \ldots, 10$, the upper (lower) signs apply when i is odd (even), c_i are the Wilson coefficients, $C_F = (N_c^2 - 1)/(2N_c)$ with $N_c = 3$, M_2 is the emitted meson, and M_1 shares the same spectator quark with the B meson. The quantities $V_i(M_2)$ account for vertex corrections, $H_i(M_1 M_2)$ 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 expression of the quantities $N_i(M_2)$ reads

$$N_i(M_2) = \begin{cases} 0, & i = 6, 8 \quad \text{and} \quad M_2 = V, \\ 1, & \text{else.} \end{cases} \tag{3.2}$$

Note that $N_i(M_2)$ vanishes for $i = 6, 8$ and $M_2 = V$ owing to the consequence of the second equation in Eq. (2.15).

The vertex and penguin corrections for SV final states have the same expressions as those for PP and PV states and can be found in [15,17]. Using the general LCDA

$$\Phi_M(x,\mu) = 6x(1-x) \left[1 + \sum_{n=1}^{\infty} \alpha_n^M(\mu) C_n^{3/2}(2x-1) \right] \tag{3.3}$$

and applying Eq. (37) in [15] for vertex corrections, we obtain

$$\begin{aligned} V_i(M) = &12 \ln \frac{m_b}{\mu} - 18 - \frac{1}{2} - 3i\pi + \left(\frac{11}{2} - 3i\pi \right) \alpha_1^M \\ &- \frac{21}{20} \alpha_2^M + \left(\frac{79}{36} - \frac{2i\pi}{3} \right) \alpha_3^M + \cdots, \end{aligned} \tag{3.4}$$

for $i = 1$–4, 9, 10

$$\begin{aligned} V_i(M) = &-12 \ln \frac{m_b}{\mu} + 6 + \frac{1}{2} + 3i\pi + \left(\frac{11}{2} - 3i\pi \right) \alpha_1^M \\ &+ \frac{21}{20} \alpha_2^M + \left(\frac{79}{36} - \frac{2i\pi}{3} \right) \alpha_3^M + \cdots, \end{aligned} \tag{3.5}$$

for $i = 5, 7$ and

$$V_i(M) = \begin{cases} -6 & \text{for } M = S, \\ 9 - 6\pi i & \text{for } M = V, \end{cases} \tag{3.6}$$

for $i = 6, 8$ in the naive dimensional regularization scheme for γ_5. The expressions of $V_i(M)$ up to the α_2^M term are the same as that in [17].

[1]For neutral scalars $f_0(980)$ and $a_0^0(980)$, it is more convenient to use the effective Wilson coefficients $\bar{a}_i$ defined as $a_i \mu_S^{-1}$; see the Appendix for more details.

As for the hard spectator function H, it reads

$$H_i(M_1M_2) = -\frac{f_Bf_{M_1}}{D(M_1M_2)}\int_0^1\frac{d\rho}{\rho}\Phi_B(\rho)\int_0^1\frac{d\xi}{\bar{\xi}}\Phi_{M_2}(\xi) \times \int_0^1\frac{d\eta}{\bar{\eta}}\left[\pm\Phi_{M_1}(\eta) + r_\chi^{M_1}\frac{\bar{\xi}}{\xi}\Phi_{m_1}(\eta)\right], \tag{3.7}$$

for $i = 1$–4, 9, 10, where the upper sign is for $M_1 = V$ and the lower sign for $M_1 = S$,

$$H_i(M_1M_2) = \frac{f_Bf_{M_1}}{D(M_1M_2)}\int_0^1\frac{d\rho}{\rho}\Phi_B(\rho)\int_0^1\frac{d\xi}{\xi}\Phi_{M_2}(\xi) \times \int_0^1\frac{d\eta}{\bar{\eta}}\left[\pm\Phi_{M_1}(\eta) + r_\chi^{M_1}\frac{\xi}{\bar{\xi}}\Phi_{m_1}(\eta)\right], \tag{3.8}$$

for $i = 5, 7$ and $H_i = 0$ for $i = 6, 8$, $\bar{\xi} \equiv 1 - \xi$ and $\bar{\eta} \equiv 1 - \eta$, Φ_M (Φ_m) is the twist-2 (twist-3) light-cone distribution amplitude of the meson M, and

$$D(SV) = F_1^{BS}(0)m_B^2, \qquad D(VS) = A_0^{BV}(0)m_B^2. \tag{3.9}$$

The ratios r_χ^V and r_χ^S are defined as

$$r_\chi^V(\mu) = \frac{2m_V}{m_b(\mu)}\frac{f_V^\perp(\mu)}{f_V},$$
$$r_\chi^S(\mu) = \frac{2m_S}{m_b(\mu)}\frac{\bar{f}_S(\mu)}{f_S} = \frac{2m_S^2}{m_b(\mu)(m_2(\mu) - m_1(\mu))}. \tag{3.10}$$

For the neutral scalars σ, f_0 and a_0^0, r_χ^S becomes divergent while f_S vanishes. In this case one needs to express $f_Sr_\chi^S$ by $\bar{f}_S\bar{r}_\chi^S$ with

$$\bar{r}_\chi^S(\mu) = \frac{2m_S}{m_b(\mu)}. \tag{3.11}$$

Weak annihilation contributions are described by the terms b_i, and $b_{i,\rm EW}$ in Eq. (A8) which have the expressions

$$\begin{aligned} b_1 &= \frac{C_F}{N_c^2}c_1A_1^i, \\ b_3 &= \frac{C_F}{N_c^2}[c_3A_1^i + c_5(A_3^i + A_3^f) + N_cc_6A_3^f], \\ b_2 &= \frac{C_F}{N_c^2}c_2A_1^i, \\ b_4 &= \frac{C_F}{N_c^2}[c_4A_1^i + c_6A_2^f], \\ b_{3,\rm EW} &= \frac{C_F}{N_c^2}[c_9A_1^i + c_7(A_3^i + A_3^f) + N_cc_8A_3^i], \\ b_{4,\rm EW} &= \frac{C_F}{N_c^2}[c_{10}A_1^i + c_8A_2^i], \end{aligned} \tag{3.12}$$

where the subscripts 1, 2, 3 of $A_n^{i,f}$ denote the annihilation amplitudes induced from $(V - A)(V - A)$, $(V - A)(V + A)$, and $(S - P)(S + P)$ operators, respectively, and the superscripts i and f refer to gluon emission from the initial and final-state quarks, respectively. Their explicit expressions can be obtained from $A_n^{i,f}(VV)$ for the VV case [18] with the replacements specified in Eq. (A1):

$$\begin{aligned}
A_1^i &= \int\cdots\begin{cases}(\Phi_V(x)\Phi_S(y)[\frac{1}{x(1-\bar{x}y)} + \frac{1}{x\bar{y}^2}] + r_\chi^Vr_\chi^S\Phi_v(x)\Phi_S^s(y)\frac{2}{x\bar{y}}); & \text{for } M_1M_2 = VS,\\ (\Phi_S(x)\Phi_V(y)[\frac{1}{x(1-\bar{x}y)} + \frac{1}{x\bar{y}^2}] + r_\chi^Vr_\chi^S\Phi_S^s(x)\Phi_v(y)\frac{2}{x\bar{y}}); & \text{for } M_1M_2 = SV,\end{cases}\\
A_2^i &= \int\cdots\begin{cases}(\Phi_V(x)\Phi_S(y)[\frac{1}{\bar{y}(1-\bar{x}y)} + \frac{1}{x^2\bar{y}}] + r_\chi^Vr_\chi^S\Phi_v(x)\Phi_S^s(y)\frac{2}{x\bar{y}}); & \text{for } M_1M_2 = VS,\\ (\Phi_S(x)\Phi_V(y)[\frac{1}{\bar{y}(1-\bar{x}y)} + \frac{1}{x^2\bar{y}}] + r_\chi^Vr_\chi^S\Phi_S^s(x)\Phi_v(y)\frac{2}{x\bar{y}}); & \text{for } M_1M_2 = SV,\end{cases}\\
A_3^i &= \int\cdots\begin{cases}(r_\chi^V\Phi_v(x)\Phi_S(y)\frac{2\bar{x}}{x\bar{y}(1-\bar{x}y)} - r_\chi^S\Phi_V(x)\Phi_S^s(y)\frac{2y}{x\bar{y}(1-\bar{x}y)}); & \text{for } M_1M_2 = VS,\\ (-r_\chi^S\Phi_S^s(x)\Phi_V(y)\frac{2\bar{x}}{x\bar{y}(1-\bar{x}y)} + r_\chi^V\Phi_S(x)\Phi_v(y)\frac{2y}{x\bar{y}(1-\bar{x}y)}); & \text{for } M_1M_2 = SV,\end{cases}\\
A_3^f &= \int\cdots\begin{cases}(r_\chi^V\Phi_v(x)\Phi_S(y)\frac{2(1+\bar{y})}{x\bar{y}^2} + r_\chi^S\Phi_V(x)\Phi_S^s(y)\frac{2(1+x)}{x^2\bar{y}}); & \text{for } M_1M_2 = VS,\\ (-r_\chi^S\Phi_S^s(x)\Phi_V(y)\frac{2(1+\bar{y})}{x\bar{y}^2} - r_\chi^V\Phi_S(x)\Phi_v(y)\frac{2(1+x)}{x^2\bar{y}}); & \text{for } M_1M_2 = SV,\end{cases} \qquad A_1^f = A_2^f = 0,
\end{aligned} \tag{3.13}$$

where $\int\cdots = \pi\alpha_s\int_0^1 dxdy$, $\bar{x} = 1 - x$, and $\bar{y} = 1 - y$. Note that we have adopted the same convention as in [15] that M_1 contains an antiquark from the weak vertex with longitudinal fraction $\bar{y}$, while M_2 contains a quark from the weak vertex with momentum fraction x.

Using the asymptotic distribution amplitudes for vector mesons and keeping the LCDA of the scalar meson to the third Gegenbauer polynomial in Eq. (2.13), the annihilation contributions can be simplified to

$$\begin{aligned}
A_1^i(VS) &\approx 6\pi\alpha_s\left\{3\mu_S\left[B_1(3X_A+4-\pi^2)+B_3\left(10X_A+\frac{23}{18}-\frac{10}{3}\pi^2\right)\right]-r_\chi^S r_\chi^V X_A(X_A-2)\right\},\\
A_2^i(VS) &\approx 6\pi\alpha_s\left\{3\mu_S\left[B_1(X_A+29-3\pi^2)+B_3\left(X_A+\frac{2956}{9}-\frac{100}{3}\pi^2\right)\right]-r_\chi^S r_\chi^V X_A(X_A-2)\right\},\\
A_3^i(VS) &\approx 6\pi\alpha_s\left\{-r_\chi^V\mu_S\left[9B_1(X_A^2-4X_A-4+\pi^2)+10B_3\left(3X_A^2-19X_A+\frac{61}{6}+3\pi^2\right)\right]-r_\chi^S\left(X_A^2-2X_A+\frac{\pi^2}{3}\right)\right\},\\
A_3^f(VS) &\approx 6\pi\alpha_s\left\{-3r_\chi^V\mu_S(X_A-2)\left[B_1(6X_A-11)+B_3\left(20X_A-\frac{187}{3}\right)\right]+r_\chi^S X_A(2X_A-1)\right\},
\end{aligned} \tag{3.14}$$

for $M_1M_2 = VS$, and

$$\begin{aligned}
A_1^i(SV) &= -A_2^i(VS), \qquad A_2^i(SV) = -A_1^i(VS),\\
A_3^i(SV) &= A_3^i(VS), \qquad A_3^f(SV) = -A_3^f(VS),
\end{aligned} \tag{3.15}$$

for $M_1M_2 = SV$, where the end point divergence X_A is defined in Eq. (3.16) below. As noticed in passing, for neutral scalars σ, f_0, and a_0^0, one needs to express $f_S r_\chi^S$ by $\bar{f}_S\bar{r}_\chi^S$ and $f_S\mu_S$ by $\bar{f}_S$. Numerically, the dominant annihilation contribution arises from the factorizable penguin-induced annihilation characterized by A_3^f. Physically, this is because the penguin-induced annihilation contribution is not subject to helicity suppression.

Although the parameters $a_i (i \neq 6, 8)$ and $a_{6,8}r_\chi$ are formally renormalization scale and γ_5 scheme independent, in practice there exists some residual scale dependence in $a_i(\mu)$ to finite order. To be specific, we shall evaluate the vertex corrections to the decay amplitude at the scale $\mu = m_b/2$. In contrast, as stressed in [17], the hard spectator and annihilation contributions should be evaluated at the hard-collinear scale $\mu_h = \sqrt{\mu\Lambda_h}$ with $\Lambda_h \approx 500$ MeV. There is one more serious complication about these contributions; that is, while QCD factorization predictions are model independent in the $m_b \to \infty$ limit, power corrections always involve troublesome end point divergences. For example, the annihilation amplitude has end point divergences even at twist-2 level and the hard spectator scattering diagram at twist-3 order is power suppressed and possess soft and collinear divergences arising from the soft spectator quark. Since the treatment of end point divergences is model dependent, subleading power corrections generally can be studied only in a phenomenological way. We shall follow [17] to parametrize the end point divergence $X_A \equiv \int_0^1 dx/\bar{x}$ in the annihilation diagram as

$$X_A = \ln\left(\frac{m_B}{\Lambda_h}\right)(1+\rho_A e^{i\phi_A}), \tag{3.16}$$

with the unknown real parameters ρ_A and ϕ_A. Likewise, the end point divergence X_H in the hard spectator contributions can be parametrized in a similar manner.

It should be stressed again that the above prescription for treating end point divergences is just a model for $1/m_b$ corrections. Besides the penguin and annihilation contributions formally of order $1/m_b$, there may exist other power corrections which in general are difficult to study as they are nonperturbative in nature. The so-called "charming penguin" contribution is one of the long-distance effects that have been widely discussed. Recently, it has been shown that such an effect can be incorporated in final-state interactions [19]. However, in order to see the relevance of the charming penguin effect to B decays into scalar resonances, we need to await more data with better accuracy.

IV. RESULTS AND DISCUSSIONS

While it is widely believed that the $f_0(980)$ and the $a_0(980)$ are predominately four-quark states, in practice it is difficult to make quantitative predictions on hadronic $B \to SV$ decays based on the four-quark picture for light scalar mesons as it involves not only the unknown form factors and decay constants that are beyond the conventional quark model but also additional nonfactorizable contributions that are difficult to estimate. Hence, we shall assume the $q\bar{q}$ scenario for the $f_0(980)$ and the $a_0(980)$ in order to apply QCDF.

For $a_0(1450)V$ and $K_0^*(1430)$ channels, we have explored in [1] two possible scenarios for the scalar mesons above 1 GeV in the QCD sum rule method: (i) In scenario 1, we treat κ, $a_0(980)$, $f_0(980)$ as the lowest lying states, and $K_0^*(1430)$, $a_0(1450)$, $f_0(1500)$ as the corresponding first excited states, respectively, where we have assumed that $f_0(980)$ and $f_0(1500)$ are dominated by the $\bar{s}s$ component and (ii) we assume in scenario 2 that $K_0^*(1430)$ $a_0(1450)$, $f_0(1500)$ are the lowest lying resonances and the corresponding first excited states lie between (2.0–2.3) GeV. Scenario 2 corresponds to the case that light scalar mesons are four-quark bound states, while all scalar mesons are made of two quarks in scenario 1. Hence, in scenario 2 we cannot make any predictions for $f_0(980)$ and $a_0(980)$ in QCD factorization.

The calculated results for the branching ratios of $B \to SV$ are shown in Tables III and IV. In these tables we have included theoretical errors arising from the uncertainties in the Gegenbauer moments $B_{1,3}$ (cf. Table II), the scalar meson decay constant f_S or $\bar{f}_S$ [see Eq. (2.6)], the form factors $F^{BP,BS}$, the quark masses and the power corrections

TABLE III. Branching ratios (in units of 10^{-6}) of B decays to final states containing a scalar meson and a vector meson. The theoretical errors correspond to the uncertainties due to (i) the Gegenbauer moments $B_{1,3}$, the scalar meson decay constants, (ii) the heavy-to-light form factors and the strange quark mass, and (iii) the power corrections due to weak annihilation and hard spectator interactions, respectively. Branching ratios of $B \to f_0(980)K^*$, $f_0(980)\rho$ are calculated for the $f_0 - \sigma$ mixing angle $\theta = 15°$. For light scalar mesons $f_0(980)$ and $a_0(980)$ we have assumed the $q\bar{q}$ content for them. The scalar mesons $a_0(1450)$ and $K_0^*(1450)$ are treated as the first excited states of $a_0(980)$ and κ, respectively, corresponding to scenario 1 as explained in Appendices B and C of [1]. Experimental results are taken from Eq. (1.1). We have assumed $\mathcal{B}(f_0(980) \to \pi^+\pi^-) = 0.50$ to obtain the experimental branching ratios for $f_0(980)V$.

Mode	Theory	Experiment	Mode	Theory	Experiment
$B^- \to f_0(980)K^{*-}$	$7.9^{+1.4+1.7+7.9}_{-1.2-1.6-3.5}$	10.4 ± 2.6	$\bar{B}^0 \to f_0(980)\bar{K}^{*0}$	$7.0^{+1.3+1.5+7.9}_{-1.2-1.4-3.4}$	$5.2 \pm 2.2 < 8.6$
$B^- \to f_0(980)\rho^-$	$0.4^{+0.0+0.1+0.0}_{-0.0-0.1-0.0}$	<3.8	$\bar{B}^0 \to f_0(980)\rho^0$	$0.03^{+0.00+0.00+0.01}_{-0.00-0.00-0.01}$	<1.06
			$\bar{B}^0 \to f_0(980)\omega$	$0.03^{+0.02+0.00+0.02}_{-0.02-0.00-0.01}$	<3.0
$B^- \to a_0^0(980)K^{*-}$	$2.8^{+0.2+0.5+6.0}_{-0.2-0.5-1.5}$		$\bar{B}^0 \to a_0^+(980)K^{*-}$	$4.5^{+0.1+0.8+3.8}_{-0.1-0.8-1.5}$	
$B^- \to a_0^-(980)\bar{K}^{*0}$	$6.1^{+0.2+1.1+4.7}_{-0.2-1.0-2.1}$		$\bar{B}^0 \to a_0^0(980)\bar{K}^{*0}$	$2.6^{+0.3+0.5+6.6}_{-0.3-0.4-1.9}$	
$B^- \to a_0^0(980)\rho^-$	$21.0^{+0.5+3.9+3.3}_{-0.5-3.6-2.3}$		$\bar{B}^0 \to a_0^+(980)\rho^-$	$27.7^{+2.2+5.5+5.1}_{-2.0-5.0-4.2}$	
$B^- \to a_0^-(980)\rho^0$	$2.7^{+0.7+0.1+2.1}_{-0.6-0.1-0.9}$		$\bar{B}^0 \to a_0^-(980)\rho^+$	$0.10^{+0.03+0.00+0.30}_{-0.02-0.00-0.07}$	
$B^- \to a_0^-(980)\omega$	$0.9^{+0.3+0.0+0.3}_{-0.2-0.0-0.3}$		$\bar{B}^0 \to a_0^0(980)\rho^0$	$1.7^{+0.3+0.2+0.2}_{-0.2-0.2-0.1}$	
			$\bar{B}^0 \to a_0^0(980)\omega$	$1.3^{+0.2+0.1+0.0}_{-0.2-0.1-0.0}$	
$B^- \to a_0^0(1450)K^{*-}$	$1.1^{+0.1+0.3+12.8}_{-0.1-0.2-0.4}$		$\bar{B}^0 \to a_0^+(1450)K^{*-}$	$0.6^{+0.0+0.3+17.4}_{-0.0-0.2-0.1}$	
$B^- \to a_0^-(1450)\bar{K}^{*0}$	$0.4^{+0.1+0.3+18.8}_{-0.1-0.2-0.2}$		$\bar{B}^0 \to a_0^0(1450)\bar{K}^{*0}$	$0.7^{+0.2+0.2+12.9}_{-0.2-0.2-0.5}$	
$B^- \to a_0^0(1450)\rho^-$	$30.5^{+1.1+3.6+6.3}_{-1.0-3.4-3.3}$		$\bar{B}^0 \to a_0^+(1450)\rho^-$	$20.7^{+3.1+4.5+17.7}_{-2.6-4.0-12.1}$	
$B^- \to a_0^-(1450)\rho^0$	$0.3^{+0.1+0.0+1.4}_{-0.0-0.0-0.2}$		$\bar{B}^0 \to a_0^-(1450)\rho^+$	$1.1^{+0.5+0.0+0.8}_{-0.3-0.0-0.8}$	
$B^- \to a_0^-(1450)\omega$	$0.4^{+0.1+0.1+0.3}_{-0.0-0.1-0.1}$		$\bar{B}^0 \to a_0^0(1450)\rho^0$	$11.9^{+2.5+0.1+7.2}_{-2.2-0.1-5.2}$	
			$\bar{B}^0 \to a_0^0(1450)\omega$	$10.2^{+2.4+0.1+5.5}_{-2.2-0.1-4.1}$	
$B^- \to K_0^{*-}(1430)\phi$	$1.0^{+0.3+0.4+20.2}_{-0.3-0.3-0.5}$		$\bar{B}^0 \to \bar{K}_0^{*0}(1430)\phi$	$0.9^{+0.3+0.4+19.3}_{-0.3-0.3-0.5}$	4.6 ± 0.9
$B^- \to \bar{K}_0^{*0}(1430)\rho^-$	$17.2^{+5.4+0.6+22.8}_{-4.4-0.6-5.6}$		$\bar{B}^0 \to K_0^{*-}(1430)\rho^+$	$12.6^{+3.6+0.4+23.1}_{-3.0-0.3-5.8}$	
$B^- \to K_0^{*-}(1430)\rho^0$	$6.2^{+2.5+0.5+4.9}_{-2.0-0.5-0.8}$		$\bar{B}^0 \to \bar{K}_0^{*0}(1430)\rho^0$	$10.0^{+2.4+0.5+12.1}_{-2.0-0.4-3.1}$	
$B^- \to K_0^{*-}(1430)\omega$	$6.1^{+1.4+0.2+9.3}_{-1.2-0.2-2.1}$		$\bar{B}^0 \to \bar{K}_0^{*0}(1430)\omega$	$6.4^{+1.4+0.3+4.0}_{-1.2-0.2-0.9}$	

from weak annihilation and hard spectator interactions characterized by the parameters X_A and X_H, respectively. For form factors we assign their uncertainties to be $\delta F^{BV,BS}(0) = \pm 0.03$; for example, $A_0^{BK^*}(0) = 0.31 \pm 0.03$ and $F_1^{BK_0^*}(0) = 0.26 \pm 0.03$. The strange quark mass is taken to be $m_s(2.1\text{ GeV}) = 90 \pm 20$ MeV. For the quantities X_A and X_H we adopt the form (3.16) with $\rho_{A,H} \leq 0.5$ and arbitrary strong phases $\phi_{A,H}$. Note that the central values (or "default" results) correspond to $\rho_{A,H} = 0$ and $\phi_{A,H} = 0$. We emphasize again that, since there are several other possible sources of power corrections as discussed by the end of Sec. III, the last errors we quote in Tables III and IV could seriously underestimate the $1/m_b$ corrections in a number of modes, especially the penguin-dominated ones.

TABLE IV. Same as Table III except that the mesons $a_0(1450)$ and $K_0^*(1450)$ are treated as the lowest lying scalar states, corresponding to scenario 2 as explained in Appendices B and C of [1].

Mode	Theory	Experiment	Mode	Theory	Experiment
$B^- \to a_0^0(1450)K^{*-}$	$2.6^{+0.2+0.4+12.7}_{-0.2-0.4-1.7}$		$\bar{B}^0 \to a_0^+(1450)K^{*-}$	$5.3^{+0.3+0.8+19.8}_{-0.3-0.8-2.6}$	
$B^- \to a_0^-(1450)\bar{K}^{*0}$	$7.8^{+0.5+1.2+23.6}_{-0.5-1.1-4.6}$		$\bar{B}^0 \to a_0^0(1450)\bar{K}^{*0}$	$2.7^{+0.4+0.5+13.8}_{-0.3-0.4-2.3}$	
$B^- \to a_0^0(1450)\rho^-$	$25.4^{+0.4+3.3+4.1}_{-0.4-3.1-3.9}$		$\bar{B}^0 \to a_0^+(1450)\rho^-$	$13.3^{+2.4+3.6+17.1}_{-2.0-3.2-8.7}$	
$B^- \to a_0^-(1450)\rho^0$	$4.5^{+1.3+0.1+6.8}_{-1.1-0.1-2.0}$		$\bar{B}^0 \to a_0^-(1450)\rho^+$	$2.6^{+1.2+0.0+1.6}_{-0.9-0.0-1.0}$	
$B^- \to a_0^-(1450)\omega$	$1.4^{+0.5+0.0+1.1}_{-0.4-0.0-0.8}$		$\bar{B}^0 \to a_0^0(1450)\rho^0$	$3.2^{+0.8+0.1+1.8}_{-0.7-0.1-1.3}$	
			$\bar{B}^0 \to a_0^0(1450)\omega$	$1.8^{+0.5+0.0+1.3}_{-0.4-0.0-0.8}$	
$B^- \to K_0^{*-}(1430)\phi$	$17.3^{+6.2+1.7+52.4}_{-4.7-1.7-12.1}$		$\bar{B}^0 \to \bar{K}_0^{*0}(1430)\phi$	$16.9^{+6.2+1.7+51.8}_{-4.7-1.6-12.0}$	4.6 ± 0.9
$B^- \to \bar{K}_0^{*0}(1430)\rho^-$	$66.0^{+24.9+2.8+71.1}_{-19.4-2.4-26.3}$		$\bar{B}^0 \to K_0^{*-}(1430)\rho^+$	$51.7^{+16.5+1.5+68.9}_{-13.4-1.4-24.0}$	
$B^- \to K_0^{*-}(1430)\rho^0$	$21.9^{+7.8+1.3+30.0}_{-6.2-1.2-10.4}$		$\bar{B}^0 \to \bar{K}_0^{*0}(1430)\rho^0$	$36.0^{+13.8+0.9+23.2}_{-10.7-0.7-9.0}$	
$B^- \to K_0^{*-}(1430)\omega$	$15.3^{+4.7+0.9+22.8}_{-3.8-0.8-7.7}$		$\bar{B}^0 \to \bar{K}_0^{*0}(1430)\omega$	$14.6^{+4.2+1.1+14.6}_{-3.5-1.0-5.1}$	

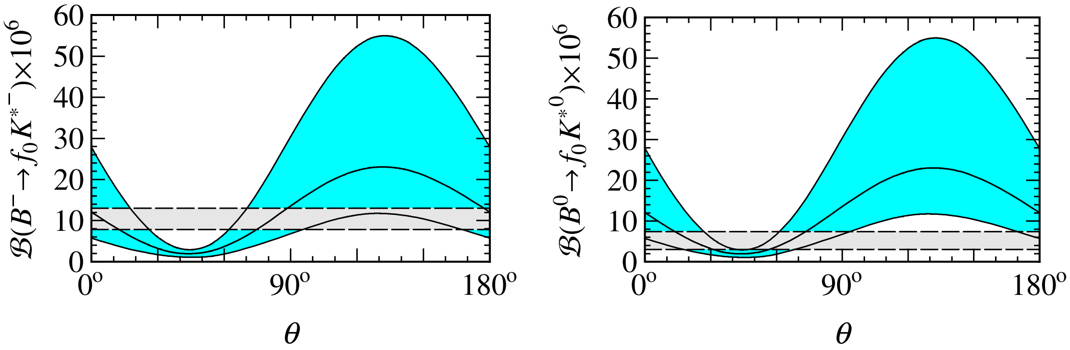

FIG. 1 (color online). Branching ratios of $B^- \to f_0(980)K^{*-}$ and $B^0 \to f_0(980)K^{*0}$ versus the mixing angle θ of strange and nonstrange components of $f_0(980)$, where the dark area is the theoretically allowed region with one sigma theoretical error and the middle bold solid curve corresponds to the central value. The horizontal band shows the experimentally allowed region with one sigma error.

A. $B \to f_0(980)K^*$ and $a_0(980, 1450)K^*$ decays

The penguin-dominated $B \to f_0(980)K^*$ decay receives three distinct types of factorizable contributions: one from the K^* emission, one from the f_0 emission with the $s\bar{s}$ content, and the other from the f_0 emission with the $n\bar{n}$ component.[2] In the expression of $B \to f_0 K^*$ decay amplitudes given in Eq. (A8), the superscript u of the form factor $F_0^{Bf_0^u}$ reminds us that it is the u quark component of f_0 that gets involved in the form factor transition. In contrast, the superscript q of the decay constant $\bar{f}_{f_0}^q$ indicates that it is the $q\bar{q}$ quark content of f_0 responsible for the penguin contribution under consideration. Note that the f_0 emission amplitude induced from four-quark operators other than O_6 and O_8 is proportional to the vanishing f_0 decay constant. However, it is compensated by the μ_S term in the twist-2 LCDA of the scalar meson so that the combination $f_{f_0}\mu_{f_0} = \bar{f}_{f_0}$ becomes finite.

In the extreme case that the $f_0(980)$ is made of $\bar{s}s$ quarks or $\bar{n}n$ quarks, the branching ratio of $B^- \to f_0(980)K^{*-}$ is given by

$$\mathcal{B}(B^- \to f_0(980)K^{*-}) = \begin{cases} (12.7^{+1.9+2.5+15.6}_{-1.7-2.3-6.1}) \times 10^{-6}; & \text{for } f_0(980) = \bar{s}s, \\ (12.5^{+3.7+1.5+15.3}_{-3.0-1.4-5.6}) \times 10^{-6}; & \text{for } f_0(980) = \bar{n}n. \end{cases} \tag{4.1}$$

In general, $\mathcal{B}(B \to f_0(980)K^*)$ depends on the mixing angle θ of strange and nonstrange components of the $f_0(980)$ (see Fig. 1). The charged and neutral modes of $f_0(980)\bar{K}^*$ are expected to have similar rates, while experimentally their central values differ by a factor of 2. This discrepancy needs to be clarified by the future improved measurements.

[2]In our previous work for $B \to SP$ decays [1], we did not take into account the contributions from the f_0 or the neutral a_0 emission induced from the four-quark operators other than O_6 and O_8 (see also [20]). Corrections will be published elsewhere.

In order to compare theory with experiment for $B \to f_0(980)K^*$, we need an input for $\mathcal{B}(f_0(980) \to \pi^+\pi^-)$. To do this, we shall use the BES measurement [21]

$$\frac{\Gamma(f_0(980) \to \pi\pi)}{\Gamma(f_0(980) \to \pi\pi) + \Gamma(f_0(980) \to K\bar{K})} = 0.75^{+0.11}_{-0.13}. \tag{4.2}$$

Assuming that the dominance of the $f_0(980)$ width by $\pi\pi$ and $K\bar{K}$ and applying isospin relation, we obtain

$$\begin{aligned} \mathcal{B}(f_0(980) \to \pi^+\pi^-) &= 0.50^{+0.07}_{-0.09}, \\ \mathcal{B}(f_0(980) \to K^+K^-) &= 0.125^{+0.018}_{-0.022}. \end{aligned} \tag{4.3}$$

Hence, we assume $\mathcal{B}(f_0(980) \to \pi^+\pi^-) = 0.50$ to determine the absolute branching ratio for $B \to f_0(980)K^*$.

For a_0K^* decays, they have similar rates as the corresponding a_0K channels [1].

B. $B \to f_0(980)\rho$ and $a_0(980\,1450)\rho$ decays

The tree-dominated decays $B \to a_0(980)\rho$, $f_0(980)\rho$ are governed by the $B \to a_0$ and $B \to f_0^u$ transition form factors, respectively. The $f_0\rho$ rate is rather small because of the small $u\bar{u}$ component in the $f_0(980)$ and the destructive interference between a_4 and a_6 penguin terms. The $f_0\rho^0$ and $f_0\omega$ modes are suppressed relative to $f_0\rho^-$ by at least a factor of $\frac{1}{2}|a_2/a_1|^2$.

The decay $\bar{B}^0 \to a_0^+\rho^-$ has a rate much larger than the $a_0^-\rho^+$ one because the factorizable amplitude of the former (latter) is proportional to f_ρ (f_{a_0}) and the decay constant of the charged a_0 is very small. We also notice that the predicted $a_0\rho^-$ rates are much larger than that of $a_0\pi^-$ for two reasons. First of all, the ρ meson decay constant is bigger than that of the pion, $f_\rho > f_\pi$. Second, the destructive interference between the a_4 and a_6 penguin terms is less severe for $a_0\rho$ as $r_\chi^\pi \sim 2.4 r_\chi^\rho$. Contrary to the naive anticipation that $\Gamma(B^- \to a_0^0\rho^-) \sim \frac{1}{2}\Gamma(\bar{B}^0 \to a_0^+\rho^-)$, we found that this relation is violated especially for $a_0 = a_0(1450)$ due to additional contributions to the $a_0^0\rho^-$ mode from the a_0^0 emission.

In general, QCD factorization works best for the color-allowed decay modes such as $a_0^0\rho^-$ and $a_0^+\rho^-$ [here $a_0 = a_0(980)$ or $a_0(1450)$] as they are tree dominated and have large branching fractions. From Tables III and IV we see that the last errors obtained for tree-dominated modes are much smaller that for penguin-dominated ones.

Recently, the isovector scalar meson $a_0(1450)$ has been confirmed to be a conventional $q\bar{q}$ meson in lattice calculations [22–26]. Hence, the calculations for the $a_0(1450)$ channels should be more trustworthy. Our results indicate that $a_0^+(1450)\rho^-$ and $a_0^0(1450)\rho^-$ have large branching ratios, of order 20×10^{-6} and 30×10^{-6}, respectively. A measurement of them at the predicted level will reinforce the $q\bar{q}$ nature for the $a_0(1450)$.

C. $B\to K_0^*(1430)\phi$ and $K_0^*(1430)\rho$ decays

For $K_0^*(1430)\phi$ channels, the central value of the predicted $\mathcal{B}(\bar{B}^0\to\bar{K}_0^*(1430)\phi)$ in scenario 1 (2) is too small (large) compared to the experimental value of $(4.6\pm0.9)\times10^{-6}$ (see Tables III and IV), though they are consistent within theoretical uncertainties. This mode was measured by *BABAR* [5] using the LASS parametrization to describe the $(K\pi)_0^{*0}$ amplitude. However, as commented in [27], while this approach is experimentally motivated, the use of the LASS parametrization is limited to the elastic region of $M(K\pi)\lesssim 2.0$ GeV, and an additional amplitude is still required for a satisfactory description of the data. Therefore, it will be interesting to see the Belle measurement for $K_0^*(1430)\phi$ modes.

Theoretically, the $K_0^*(1430)\rho$ rates are expected to be substantially larger than that of the $K_0^*(1430)\pi$ ones since the penguins terms a_4 and a_6 contribute constructively to the former and destructively to the latter. However, as shown in [1], our predicted central values for the branching ratios of $\bar{K}_0^{*0}\pi^-$ and $K_0^{*-}\pi^+$ are too small by a factor 3–4 compared to experiment.[3] It appears that one needs sizable weak annihilation in order to accommodate the $K_0^*\pi$ data. In this work, we found large rates for $\bar{K}_0^{*0}\rho^{-,0}$ and $K_0^{*-}\rho^+$ even in the absence of weak annihilation contributions. Experimentally, it should be relatively easy to search for those $K_0^*(1430)\rho$ modes to see if they are enhanced relative to their counterparts in the $K_0^*\pi$ sector. The branching ratios for the $K_0^*(1430)\omega$ modes are predicted to be of order $(6\sim15)\times10^{-6}$.

[3]Recently, the authors of [28] claimed that the decay rates for the $\bar{K}_0^{*0}\pi^-$ and $K_0^{*-}\pi^+$ modes can be accommodated in the pQCD approach. It is not clear to us what is the underlying reason for the discrepancy between our work and [28]. However, we have just performed a systematical study of charmless 3-body B decays based on a simple generalized factorization approach [29]. We consider the weak process $B\to K_0^*(1430)\pi$ followed by the strong decay $K_0^*\to K\pi$ and reach the same conclusion as [1], namely, the predicted $\bar{K}_0^{*0}\pi^-$ and $K_0^{*-}\pi^+$ rates are too small compared to the data.

V. CONCLUSIONS

We have studied the hadronic charmless B decays into a scalar meson and a vector meson within the framework of QCD factorization. The main results are:

(i) The decay rates for the $f_0(980)K^{*-}$ and $f_0(980)\bar{K}^{*0}$ modes depend on the mixing angle θ of strange and nonstrange components of the $f_0(980)$.

(ii) QCD factorization works best for the color-allowed decay modes such as $a_0^0\rho^-$ and $a_0^+\rho^-$ [$a_0 = a_0(980)$ or $a_0(1450)$] as they are tree dominated and have large branching fractions.

(iii) If the $a_0(980)$ is a $q\bar{q}$ bound state, the predicted branching ratios for the channels $a_0^+\rho^-$ and $a_0^0\rho^-$ will be very large, of order 28×10^{-6} and 21×10^{-6}, respectively.

(iv) For the $a_0(1450)$ channels, $a_0^+(1450)\rho^-$ and $a_0^0(1450)\rho^-$ are found to have branching ratios of order 20×10^{-6} and 30×10^{-6}, respectively. An observation of them at the predicted level will favor the $q\bar{q}$ structure for the $a_0(1450)$.

(v) Contrary to the naive expectation that $\Gamma(B^-\to a_0^0\rho^-)\sim\frac{1}{2}\Gamma(\bar{B}^0\to a_0^+\rho^-)$, we found that this naive relation is violated especially for $a_0 = a_0(1450)$ due to additional contributions to the $a_0^0\rho^-$ mode from the a_0^0 emission.

(vi) The decays $B\to K_0^*(1430)\rho$ are expected to have rates substantially larger than that of $B\to K_0^*(1430)\pi$ owing to the constructive (destructive) interference between the a_4 and a_6 penguin terms in the former (latter). Experimentally, it is thus important to check if the $B\to K_0^*\rho$ modes are enhanced relative to their counterparts in the $K_0^*\pi$ sector. The branching ratios for the $K_0^*(1430)\omega$ modes are predicted to be of order $(6\sim15)\times10^{-6}$.

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APPENDIX A

The $B\to SV$ (VS) decay amplitudes can be either evaluated directly or obtained readily from $B\to VV$ amplitudes with the replacements:

$$\begin{aligned} &\Phi_V(x)\to\Phi_S(x), \qquad \Phi_v(x)\to\Phi_S^s(x),\\ &f_V\to f_S, \qquad f_V^\perp\to-\bar{f}_S, \qquad r_\chi^V\to-r_\chi^S. \end{aligned} \tag{A1}$$

As stressed in the main text, we use the LCDAs with the decay constants being factored out. Since the VV channels have been studied in detail in [18], we may use them to obtain the $B \to SV$ amplitudes. In [18], the factorizable longitudinal $B \to VV$ amplitude reads (apart from the effective Wilson coefficients)

$$A_{V_1V_2} = i\frac{G_F}{\sqrt{2}} 2f_{V_2} A_0^{BV_1}(m_{V_2}^2) m_B p_c, \tag{A2}$$

where use has been made of the replacement $m_V \varepsilon^* \cdot p_B \to m_B p_c$ with p_c being the c.m. momentum. Since the definitions for the decay constant f_V and the form factor A_0 in [18]

$$\begin{aligned} \langle V|V_\mu|0\rangle &= -if_V m_V \varepsilon_\mu^*, \\ \langle V(p')|A_\mu|B(p)\rangle &= 2m_V \frac{\varepsilon^* \cdot P}{q^2} q_\mu A_0^{PV}(q^2) + \cdots \end{aligned} \tag{A3}$$

are different from ours [see Eqs. (2.1) and (2.8)], the replacements (A1) need to be modified accordingly. The $B \to VS$ amplitude is obtained from the replacements:

$$\begin{aligned} &f_{V_1} \to if_V, \qquad f_{V_2} \to if_S, \\ &A_0^{BV_1} \to iA_0^{BV}, \qquad r_\chi^{V_2} \to -r_\chi^S. \end{aligned} \tag{A4}$$

For $B \to SV$ amplitudes, the replacements are

$$\begin{aligned} &f_{V_1} \to if_S, \qquad f_{V_2} \to if_V, \\ &A_0^{BV_1} \to -iF_1^{BS}, \qquad r_\chi^{V_2} \to r_\chi^V. \end{aligned} \tag{A5}$$

From (A4) and (A5) we obtain the factorizable $B \to SV$ and VS amplitudes

$$A_{M_1M_2} = i\frac{G_F}{\sqrt{2}} \begin{cases} 2f_V F_1^{BS}(m_V^2) m_B p_c; & \text{for } M_1M_2 = SV, \\ -2f_S A_0^{BV}(m_S^2) m_B p_c; & \text{for } M_1M_2 = VS. \end{cases} \tag{A6}$$

The coefficients of the flavor operators α_i^p for SV can be obtained from the VV case [18] and they read

$$\begin{aligned} &\alpha_1(M_1, M_2) = a_1(M_1, M_2), \qquad \alpha_2(M_1M_2) = a_2(M_1M_2), \qquad \alpha_3^p(M_1M_2) = a_3^p(M_1M_2) + a_5^p(M_1M_2), \\ &\alpha_{3,\rm EW}^p(M_1M_2) = a_9^p(M_1M_2) + a_7^p(M_1M_2), \qquad \alpha_4^p(M_1M_2) = \begin{cases} a_4^p(M_1M_2) - r_\chi^V a_6^p(M_1M_2); & \text{for } M_1M_2 = SV, \\ a_4^p(M_1M_2) + r_\chi^S a_6^p(M_1M_2); & \text{for } M_1M_2 = VS, \end{cases} \\ &\alpha_{4,\rm EW}^p(M_1M_2) = \begin{cases} a_{10}^p(M_1M_2) - r_\chi^V a_8^p(M_1M_2); & \text{for } M_1M_2 = SV, \\ a_{10}^p(M_1M_2) + r_\chi^S a_8^p(M_1M_2); & \text{for } M_1M_2 = VS. \end{cases} \end{aligned} \tag{A7}$$

Applying the replacement (A1)–(A7) to the $B \to VV$ amplitudes in [18], we obtain the following the factorizable amplitudes of the decays $B \to (f_0, a_0)K^*$, $f_0(\rho, \omega)$, $a_0(\rho, \omega)$, $a_0 K^*$, $K_0^*(\phi, \rho, \omega)$:

$$\begin{aligned} A(B^- \to f_0 K^{*-}) = i\frac{G_F}{\sqrt{2}} \sum_{p=u,c} \lambda_p^{(s)} \Big\{ & (a_1\delta_u^p + a_4^p - r_\chi^{K^*}(a_6^p + a_8^p) + a_{10}^p)_{f_0^u K^*} 2f_{K^*} F_1^{Bf_0^u}(m_{K^*}^2) m_B p_c \\ & - \left(\bar a_3 + \bar a_4^p + \bar a_5 + \left(a_6^p - \frac{1}{2}a_8^p\right)\bar r_\chi^{f_0} - \frac{1}{2}(\bar a_7 + \bar a_9 + \bar a_{10}^p)\right)_{K^* f_0^s} 2\bar f_{f_0}^s A_0^{BK^*}(m_{f_0}^2) m_B p_c \\ & - \left(\bar a_2\delta_u^p + 2(\bar a_3 + \bar a_5) + \frac{1}{2}(\bar a_7 + \bar a_9)\right)_{K^* f_0^u} 2\bar f_{f_0}^u A_0^{BK^*}(m_{f_0}^2) m_B p_c - f_B f_{K^*}[\bar f_{f_0^u}(\bar b_2\delta_u^p + \bar b_3 + \bar b_{3,\rm EW})_{f_0^u K^*} \\ & + \bar f_{f_0^s}(\bar b_2\delta_u^p + \bar b_3 + \bar b_{3,\rm EW})_{K^* f_0^s}]\Big\}, \\ A(\bar B^0 \to f_0 \bar K^{*0}) = i\frac{G_F}{\sqrt{2}} \sum_{p=u,c} \lambda_p^{(s)} \Big\{ & \left(a_4^p - r_\chi^{K^*}\left(a_6^p - \frac{1}{2}a_8^p\right) - \frac{1}{2}a_{10}^p\right)_{f_0^d K^*} 2f_{K^*} F_1^{Bf_0^u}(m_{K^*}^2) m_B p_c \\ & - \left(\bar a_3 + \bar a_4^p + \bar a_5 + \left(a_6^p - \frac{1}{2}a_8^p\right)\bar r_\chi^{f_0} - \frac{1}{2}(\bar a_7 + \bar a_9 + \bar a_{10}^p)\right)_{K^* f_0^s} 2\bar f_{f_0}^s A_0^{BK^*}(m_{f_0}^2) m_B p_c \\ & - \left(\bar a_2\delta_u^p + 2(\bar a_3 + \bar a_5) + \frac{1}{2}(\bar a_7 + \bar a_9)\right)_{K^* f_0^u} 2\bar f_{f_0}^u A_0^{BK^*}(m_{f_0}^2) m_B p_c \\ & - f_B f_{K^*}\left[\bar f_{f_0^d}\left(\bar b_3 - \frac{1}{2}\bar b_{3,\rm EW}\right)_{f_0^d K^*} + f_{f_0^s}\left(\bar b_3 - \frac{1}{2}\bar b_{3,\rm EW}\right)_{K^* f_0^s}\right]\Big\}, \end{aligned}$$

$$A(B^- \to a_0^0 K^{*-}) = i\frac{G_F}{2}\sum_{p=u,c}\lambda_p^{(s)}\{(a_1\delta_u^p + a_4^p - r_\chi^{K^*}(a_6^p + a_8^p) + a_{10}^p)_{a_0K^*}2f_{K^*}F_1^{Ba_0}(m_{K^*}^2)m_Bp_c - (\bar{a}_2\delta_u^p)_{K^*a_0}2\bar{f}_{a_0}A_0^{BK^*}(m_{a_0}^2)m_Bp_c - f_Bf_{K^*}\bar{f}_{a_0}(\bar{b}_2\delta_u^p + \bar{b}_3 + \bar{b}_{3,\mathrm{EW}})_{a_0K^*}\},$$

$$A(B^- \to a_0^- \bar{K}^{*0}) = i\frac{G_F}{\sqrt{2}}\sum_{p=u,c}\lambda_p^{(s)}\left\{\left(a_4^p - r_\chi^{K^*}\left(a_6^p - \frac{1}{2}a_8^p\right) - \frac{1}{2}a_{10}^p\right)_{a_0K^*}2f_{K^*}F_1^{Ba_0}(m_{K^*}^2)m_Bp_c - f_Bf_{K^*}f_{a_0}(b_2\delta_u^p + b_3 + b_{3,\mathrm{EW}})_{a_0K^*}\right\},$$

$$A(\bar{B}^0 \to a_0^+ K^{*-}) = i\frac{G_F}{\sqrt{2}}\sum_{p=u,c}\lambda_p^{(s)}\left\{\left(a_1\delta_u^p + a_4^p - r_\chi^{K^*}(a_6^p + a_8^p) + a_{10}^p)_{a_0K^*}2f_{K^*}F_1^{Ba_0}(m_{K^*}^2)m_Bp_c - f_Bf_{K^*}f_{a_0}\left(b_3 - \frac{1}{2}b_{3,\mathrm{EW}}\right)_{a_0K^*}\right]\right\},$$

$$A(\bar{B}^0 \to a_0^0 \bar{K}^{*0}) = i\frac{G_F}{2}\sum_{p=u,c}\lambda_p^{(s)}\left\{-\left(a_4^p - r_\chi^{K^*}\left(a_6^p - \frac{1}{2}a_8^p\right) - \frac{1}{2}a_{10}^p\right)_{a_0K^*}2f_{K^*}F_1^{Ba_0}(m_{K^*}^2)m_Bp_c - (\bar{a}_2\delta_u^p)_{K^*a_0}2\bar{f}_{a_0}A_0^{BK^*}(m_{a_0}^2)m_Bp_c - f_Bf_{K^*}\bar{f}_{a_0}\left(-\bar{b}_3 + \frac{1}{2}\bar{b}_{3,\mathrm{EW}}\right)_{a_0K^*}\right\},$$

$$A(B^- \to f_0\rho^-) = i\frac{G_F}{\sqrt{2}}\sum_{p=u,c}\lambda_p^{(d)}\left\{(a_1\delta_u^p + a_4^p - r_\chi^\rho(a_6^p + a_8^p) + a_{10}^p)_{f_0^u\rho}2f_\rho F_1^{Bf_0^u}(m_\rho^2)m_Bp_c - (\bar{a}_2\delta_u^p + 2(\bar{a}_3 + \bar{a}_5) + \bar{a}_4 + \left(a_6^p - \frac{1}{2}a_8^p\right)\bar{r}_\chi^{f_0} + \frac{1}{2}(\bar{a}_7 + \bar{a}_9 - \bar{a}_{10}))_{\rho f_0^u}2\bar{f}_{f_0}^u A_0^{B\rho}(m_{f_0}^2)m_Bp_c - f_Bf_\rho\bar{f}_{f_0}^u[(\bar{b}_2\delta_u^p + \bar{b}_3 + \bar{b}_{3,\mathrm{EW}})_{f_0^u\rho} + (\bar{b}_2\delta_u^p + \bar{b}_3 + \bar{b}_{3,\mathrm{EW}})_{\rho f_0^u}]\right\},$$

$$A(\bar{B}^0 \to f_0\rho^0) = i\frac{G_F}{2}\sum_{p=u,c}\lambda_p^{(d)}\left\{\left(a_2\delta_u^p - a_4^p + r_\chi^\rho\left(a_6^p - \frac{1}{2}a_8^p\right) + \frac{3}{2}(a_9^p + a_7^p) + \frac{1}{2}a_{10}^p\right)_{f_0^d\rho}2f_\rho F_1^{Bf_0^d}(m_\rho^2)m_Bp_c + \left(\bar{a}_2\delta_u^p + 2(\bar{a}_3 + \bar{a}_5) + \bar{a}_4 + \left(a_6^p - \frac{1}{2}a_8^p\right)\bar{r}_\chi^{f_0} + \frac{1}{2}(\bar{a}_7 + \bar{a}_9 - \bar{a}_{10})\right)_{\rho f_0^u}2\bar{f}_{f_0}^u A_0^{B\rho}(m_{f_0}^2)m_Bp_c - f_Bf_\rho\bar{f}_{f_0}^u[(\bar{b}_1\delta_u^p - \bar{b}_3 + \frac{1}{2}\bar{b}_{3,\mathrm{EW}} + \frac{3}{2}\bar{b}_{4,\mathrm{EW}})_{f_0^d\rho} + (\bar{b}_1\delta_u^p - \bar{b}_3 + \frac{1}{2}\bar{b}_{3,\mathrm{EW}} + \frac{3}{2}\bar{b}_{4,\mathrm{EW}})_{\rho f_0^d}]\right\},$$

$$A(\bar{B}^0 \to f_0\omega) = i\frac{G_F}{2}\sum_{p=u,c}\lambda_p^{(d)}\left\{\left(a_2\delta_u^p + a_4^p - r_\chi^\omega\left(a_6^p - \frac{1}{2}a_8^p\right) + \frac{1}{2}(a_9^p + a_7^p) - \frac{1}{2}a_{10}^p\right)_{f_0^d\omega}2f_\omega F_1^{Bf_0^d}(m_\omega^2)m_Bp_c - \left(\bar{a}_2\delta_u^p + 2(\bar{a}_3 + \bar{a}_5) + \bar{a}_4 + \left(a_6^p - \frac{1}{2}a_8^p\right)\bar{r}_\chi^{f_0} + \frac{1}{2}(\bar{a}_7 + \bar{a}_9 - \bar{a}_{10})\right)_{\omega f_0^d}2\bar{f}_{f_0}^d A_0^{B\omega}(m_{f_0}^2)m_Bp_c - f_Bf_\omega\bar{f}_{f_0}^d\left[\left(\bar{b}_1\delta_u^p + \bar{b}_3 + 2\bar{b}_4 - \frac{1}{2}\bar{b}_{3,\mathrm{EW}} + \frac{1}{2}\bar{b}_{4,\mathrm{EW}}\right)_{f_0^d\omega} + \left(\bar{b}_1\delta_u^p + \bar{b}_3 + 2\bar{b}_4 - \frac{1}{2}\bar{b}_{3,\mathrm{EW}} + \frac{1}{2}\bar{b}_{4,\mathrm{EW}}\right)_{\omega f_0^d}\right]\right\},$$

$$A(\bar{B}^0 \to a_0^+\rho^-) = i\frac{G_F}{\sqrt{2}}\sum_{p=u,c}\lambda_p^{(d)}\left\{(a_1\delta_u^p + a_4^p - r_\chi^\rho(a_6^p + a_8^p) + a_{10}^p)_{a_0\rho}2f_\rho F_1^{Ba_0}(m_\rho^2)m_Bp_c - f_Bf_\rho f_{a_0}\left[\left(b_3 + b_4 - \frac{1}{2}b_{3,\mathrm{EW}} - \frac{1}{2}b_{4,\mathrm{EW}}\right)_{a_0\rho} + (b_1\delta_u^p + b_4 + b_{4,\mathrm{EW}})_{\rho a_0}\right]\right\},$$

$$A(\bar{B}^0 \to a_0^- \rho^+) = i\frac{G_F}{\sqrt{2}}\sum_{p=u,c}\lambda_p^{(d)}\Big\{-(a_1\delta_u^p + a_4^p + r_\chi^{a_0}(a_6^p + a_8^p) + a_{10}^p)_{\rho a_0} 2 f_{a_0} A_0^{B\rho}(m_\rho^2) m_B p_c$$
$$- f_B f_\rho f_{a_0}\Big[\Big(b_3 + b_4 - \frac{1}{2}b_{3,\mathrm{EW}} - \frac{1}{2}b_{4,\mathrm{EW}}\Big)_{\rho a_0} + (b_1\delta_u^p + b_4 + b_{4,\mathrm{EW}})_{a_0\rho}\Big]\Big\},$$

$$A(B^- \to a_0^0 \rho^-) = i\frac{G_F}{2}\sum_{p=u,c}\lambda_p^{(d)}\Big\{(a_1\delta_u^p + a_4^p - r_\chi^{\rho}(a_6^p + a_8^p) + a_{10}^p)_{a_0\rho} 2 f_\rho F_1^{Ba_0}(m_\rho^2) m_B p_c$$
$$- \Big(\bar{a}_2\delta_u^p - \bar{a}_4 - \Big(a_6^p - \frac{1}{2}a_8^p\Big)\bar{r}_\chi^{a_0} + \frac{3}{2}(\bar{a}_7 + \bar{a}_9) + \frac{1}{2}\bar{a}_{10}\Big)_{\rho a_0} 2\bar{f}_{a_0} A_0^{B\rho}(m_{a_0}^2) m_B p_c$$
$$- f_B f_\rho \bar{f}_{a_0}[(\bar{b}_2\delta_\mu^p + \bar{b}_3 + \bar{b}_{3,\mathrm{EW}})_{a_0\rho} - (\bar{b}_2\delta_\mu^p + \bar{b}_3 + \bar{b}_{3,\mathrm{EW}})_{\rho a_0}]\Big\},$$

$$A(B^- \to a_0^- \rho^0) = i\frac{G_F}{2}\sum_{p=u,c}\lambda_p^{(d)}\Big\{-(a_1\delta_u^p + a_4^p + r_\chi^{a_0}(a_6^p + a_8^p) + a_{10}^p)_{\rho a_0} 2 f_{a_0} A_0^{B\rho}(m_{a_0}^2) m_B p_c$$
$$+ \Big[a_2\delta_u^p - a_4^p + r_\chi^{\rho}\Big(a_6^p - \frac{1}{2}a_8^p\Big) + \frac{1}{2}a_{10}^p + \frac{3}{2}(a_9 + a_7)\Big]_{a_0\rho} 2 f_\rho F_1^{Ba_0}(m_\rho^2) m_B p_c$$
$$- f_B f_\rho f_{a_0}[(b_2\delta_\mu^p + b_3 + b_{3,\mathrm{EW}})_{\rho a_0} - (b_2\delta_\mu^p + b_3 + b_{3,\mathrm{EW}})_{a_0\rho}]\Big\},$$

$$A(\bar{B}^0 \to a_0^0 \rho^0) = -i\frac{G_F}{2\sqrt{2}}\sum_{p=u,c}\lambda_p^{(d)}\Big\{\Big(a_2\delta_u^p - a_4^p + r_\chi^{\rho}\Big(a_6^p - \frac{1}{2}a_8^p\Big) + \frac{3}{2}(a_9 + a_7) + \frac{1}{2}a_{10}^p\Big)_{a_0\rho} 2 f_\rho F_1^{Ba_0}(m_\rho^2) m_B p_c$$
$$- \Big(\bar{a}_2\delta_u^p - \bar{a}_4 - \Big(a_6^p - \frac{1}{2}a_8^p\Big)\bar{r}_\chi^{a_0} + \frac{3}{2}(\bar{a}_7 + \bar{a}_9) + \frac{1}{2}\bar{a}_{10}\Big)_{\rho a_0} 2\bar{f}_{a_0} A_0^{B\rho}(m_{a_0}^2) m_B p_c$$
$$+ f_B f_\rho \bar{f}_{a_0}\Big[\Big(\bar{b}_1\delta_u^p + \bar{b}_3 + 2\bar{b}_4 - \frac{1}{2}(\bar{b}_{3,\mathrm{EW}} - \bar{b}_{4,\mathrm{EW}})\Big)_{a_0\rho} + \Big(\bar{b}_1\delta_u^p + \bar{b}_3 + 2\bar{b}_4 - \frac{1}{2}(\bar{b}_{3,\mathrm{EW}} - \bar{b}_{4,\mathrm{EW}})\Big)_{\rho a_0}\Big]\Big\},$$

$$A(B^- \to a_0^- \omega) = i\frac{G_F}{2}\sum_{p=u,c}\lambda_p^{(d)}\Big\{-(a_1\delta_u^p + a_4^p + r_\chi^{a_0}(a_6^p + a_8^p) + a_{10}^p)_{\omega a_0} 2 f_{a_0^-} A_0^{B\omega}(m_{a_0}^2) m_B p_c$$
$$+ \Big[a_2\delta_u^p + 2(a_3 + a_5) + a_4^p - r_\chi^{\omega}\Big(a_6^p - \frac{1}{2}a_8^p\Big) - \frac{1}{2}a_{10}^p + \frac{1}{2}(a_9 + a_7)\Big]_{a_0\omega} 2 f_\omega F_1^{Ba_0}(m_\omega^2) m_B p_c$$
$$- f_B f_\omega f_{a_0}[(b_2\delta_\mu^p + b_3 + b_{3,\mathrm{EW}})_{\omega a_0} + (b_2\delta_\mu^p + b_3 + b_{3,\mathrm{EW}})_{a_0\omega}]\Big\},$$

$$A(\bar{B}^0 \to a_0^0 \omega) = i\frac{G_F}{2\sqrt{2}}\sum_{p=u,c}\lambda_p^{(d)}\Big\{\Big(a2\delta_u^p + 2(a_3 + a_5) + a_4^p - r_\chi^{\omega}\Big(a_6^p - \frac{1}{2}a_8^p\Big) - \frac{1}{2}a_{10}^p + \frac{1}{2}(a_9 + a_7)\Big)_{a_0\omega} 2 f_\omega F_1^{Ba_0}(m_\omega^2) m_B p_c$$
$$- \Big(\bar{a}_2\delta_u^p + 2(\bar{a}_3 + \bar{a}_5) + \bar{a}_4 + \Big(a_6^p - \frac{1}{2}a_8^p\Big)\bar{r}_\chi^{a_0} + \frac{1}{2}(\bar{a}_7 + \bar{a}_9 - \bar{a}_{10})\Big)_{\omega a_0} 2\bar{f}_{a_0} A_0^{B\omega}(m_{a_0}^2) m_B p_c$$
$$- f_B f_\omega \bar{f}_{a_0}\Big[\Big(-\bar{b}_1\delta_\mu^p + \bar{b}_3 - \frac{1}{2}\bar{b}_{3,\mathrm{EW}} - \frac{3}{2}\bar{b}_{4,\mathrm{EW}}\Big)_{a_0\omega} + \Big(-\bar{b}_1\delta_\mu^p + \bar{b}_3 - \frac{1}{2}\bar{b}_{3,\mathrm{EW}} - \frac{3}{2}\bar{b}_{4,\mathrm{EW}}\Big)_{\omega a_0}\Big]\Big\},$$

$$A(B^- \to K_0^{*-}\phi) = i\frac{G_F}{\sqrt{2}}\sum_{p=u,c}\lambda_p^{(s)}\Big\{\Big(a_3 + a_4^p + a_5 - r_\chi^{\phi}\Big(a_6^p - \frac{1}{2}a_8^p\Big) - \frac{1}{2}(a_7 + a_9 + a_{10}^p)\Big)_{K_0^*\phi} 2 f_\phi F_1^{BK_0^*}(m_\phi^2) m_B p_c$$
$$- f_B f_\phi f_{K_0^*}(b_2\delta_u^p + b_3 + b_{3,\mathrm{EW}})_{K_0^*\phi}\Big\},$$

$$A(\bar{B}^0 \to \bar{K}_0^{*0}\phi) = i\frac{G_F}{\sqrt{2}}\sum_{p=u,c}\lambda_p^{(s)}\Big\{\Big(a_3 + a_4^p + a_5 - r_\chi^{\phi}\Big(a_6^p - \frac{1}{2}a_8^p\Big) - \frac{1}{2}(a_7 + a_9 + a_{10}^p)\Big)_{K_0^*\phi} 2 f_\phi F_1^{BK_0^*}(m_\phi^2) m_B p_c$$
$$- f_B f_\phi f_{K_0^*}\Big(b_3 - \frac{1}{2}b_{3,\mathrm{EW}}\Big)_{K_0^*\phi}\Big\},$$

$$A(B^-\to\bar K_0^{*0}\rho^-)=i\frac{G_F}{\sqrt2}\sum_{p=u,c}\lambda_p^{(s)}\Big\{-\Big(a_4^p+r_\chi^{K_0^*}\Big(a_6^p-\frac12a_8^p\Big)-\frac12a_{10}^p\Big)_{\rho K_0^*}2f_{K_0^*}A_0^{B\rho}(m_{K_0^*}^2)m_Bp_c$$
$$-f_Bf_\rho f_{K_0^*}\Big(b_2\delta_u^p+b_3+b_{3,\mathrm{EW}}\Big)_{\rho K_0^*}\Big\},$$
$$A(B^-\to K_0^{*-}\rho^0)=i\frac{G_F}{2}\sum_{p=u,c}\lambda_p^{(s)}\Big\{-(a_1\delta_u^p+a_4^p+r_\chi^{K_0^*}(a_6^p+a_8^p)+a_{10}^p)_{\rho K_0^*}2f_{K_0^*}A_0^{B\rho}(m_{K_0^*}^2)m_Bp_c$$
$$+\Big[a_2\delta_u^p+\frac32(a_9+a_7)\Big]_{K_0^*\rho}2f_\rho F_1^{BK_0^*}(m_\rho^2)m_Bp_c-f_Bf_\rho f_{K_0^*}(b_2\delta_u^p+b_3+b_{3,\mathrm{EW}})_{\rho K_0^*}\Big\},$$
$$A(\bar B^0\to K_0^{*-}\rho^+)=i\frac{G_F}{\sqrt2}\sum_{p=u,c}\lambda_p^{(s)}\Big\{-(a_1\delta_u^p+a_4^p+r_\chi^{K_0^*}a_6^p+a_{10}^p+r_\chi^{K_0^*}a_8^p)_{\rho K_0^*}2f_{K_0^*}A_0^{B\rho}(m_{K_0^*}^2)m_Bp_c$$
$$-f_Bf_\rho f_{K_0^*}(b_3-\frac12b_{3,\mathrm{EW}})_{\rho K_0^*}\Big\},$$
$$A(\bar B^0\to\bar K_0^{*0}\rho^0)=i\frac{G_F}{2}\sum_{p=u,c}\lambda_p^{(s)}\Big\{-\Big(-a_4^p-r_\chi^{K_0^*}\Big(a_6^p-\frac12a_8^p\Big)+\frac12a_{10}^p\Big)_{\rho K_0^*}2f_{K_0^*}A_0^{B\rho}(m_{K_0^*}^2)m_Bp_c$$
$$+\Big[a_2\delta_u^p+\frac32(a_9+a_7)\Big]_{K_0^*\rho}2f_\rho F_1^{BK_0^*}(m_\rho^2)m_Bp_c-f_Bf_\rho f_{K_0^*}\Big(-b_3+\frac12b_{3,\mathrm{EW}}\Big)_{\rho K_0^*}\Big\},$$
$$A(B^-\to K_0^{*-}\omega)=i\frac{G_F}{2}\sum_{p=u,c}\lambda_p^{(s)}\Big\{\Big[a_2\delta_u^p+2(a_3+a_5)+\frac12(a_9+a_7)\Big]_{K_0^*\omega}2f_\omega F_1^{BK_0^*}(m_\omega^2)m_Bp_c$$
$$-(a_1\delta_u^p+a_4^p+r_\chi^{K_0^*}(a_6^p+a_8^p)+a_{10}^p)_{\omega K_0^*}2f_{K_0^*}A_0^{B_\omega}(m_{K_0^*}^2)m_Bp_c-f_Bf_\omega f_{K_0^*}(b_2\delta_u^p+b_3+b_{3,\mathrm{EW}})_{\omega K_0^*}\Big\},$$
$$A(\bar B^0\to\bar K_0^{*0}\omega)=i\frac{G_F}{2}\sum_{p=u,c}\lambda_p^{(s)}\Big\{\Big[a_2\delta_u^p+2(a_3+a_5)+\frac12(a_9+a_7)\Big]_{K_0^*\omega}2f_\omega F_1^{BK_0^*}(m_\omega^2)m_Bp_c$$
$$-\Big(a_4^p+r_\chi^{K_0^*}\Big(a_6^p-\frac12a_8^p\Big)-\frac12a_{10}^p\Big)_{\omega K_0^*}2f_{K_0^*}A_0^{B\rho}(m_{K_0^*}^2)m_Bp_c-f_Bf_\omega f_{K_0^*}\Big(b_3-\frac12b_{3,\mathrm{EW}}\Big)_{\omega K_0^*}\Big\},\qquad\text{(A8)}$$

where $\lambda_p^{(q)}\equiv V_{pb}V_{pq}^*$ with $q=d$, s and

$$\bar r_\chi^{f_0}(\mu)=\frac{2m_{f_0}}{m_b(\mu)},\qquad\bar r_\chi^{a_0^0}(\mu)=\frac{2m_{a_0^0}}{m_b(\mu)},$$
$$r_\chi^{a_0^\pm}(\mu)=\frac{2m_{a_0^\pm}^2}{m_b(\mu)(m_d(\mu)-m_u(\mu))},\qquad\text{(A9)}$$

In Eq. (A8), we encounter terms such as $a_if_{f_0}$, which appears to vanish at first sight as $f_{f_0}=0$. However, when f_{f_0} combines with μ_{f_0} appearing in the twist-2 LCDA of the scalar meson [see Eq. (2.13)], it becomes finite, namely, $f_{f_0}\mu_{f_0}=\bar f_{f_0}$. Therefore, the effective Wilson coefficients $\bar a_i$ in Eq. (A8) are defined as $a_i\mu_S^{-1}$ and they can be obtained from Eq. (3.1) by retaining only those terms that are proportional to μ_S. Specifically,

$$\bar a_i^p(M_1M_2)=\frac{c_{i\pm1}}{N_c}\frac{C_F\alpha_s}{4\pi}\Big[\bar V_i(M_2)+\frac{4\pi^2}{N_c}\bar H_i(M_1M_2)\Big]+\bar P_i^p(M_2).\qquad\text{(A10)}$$

The LCDA of the neutral scalar meson in the bar quantities, $\bar V_i(S)$, $\bar P_i(S)$, and $\bar H_i(M1,M2)$ is replaced by $\bar\Phi_S$ which has the similar expression as Eq. (2.13) except that the first constant term does not contribute and the term $f_S\mu_S$ is factored out:

$$\bar\Phi_S(x,\mu)=6x(1-x)\sum_{m=1}^\infty B_m(\mu)C_m^{3/2}(2x-1).\qquad\text{(A11)}$$

In Eq. (A10),

$$\bar V_i(S)=\begin{cases}(\frac{11}{2}-3i\pi)B_1^S+(\frac{79}{36}-\frac{2i\pi}{3})B_3^S+\cdots; & \text{for } i=1\text{–}4,5,7,9,10,\\ 0; & \text{for } i=6,8.\end{cases}\qquad\text{(A12)}$$

The annihilation terms $\bar b_i$ have the same expressions as Eq. (3.12) with r_χ^S and μ_SB_i replaced by $\bar r_\chi^S$ and B_i, respectively.

For the Cabibbo-Kobayashi-Maskawa matrix elements, we use the Wolfenstein parameters $A=0.818$, $\lambda=0.225\,68$, $\bar\rho=0.141$, and $\bar\eta=0.348$ [30]. For the running quark masses we shall use

$$m_b(m_b) = 4.2\ \text{GeV}, \quad m_b(2.1\ \text{GeV}) = 4.95\ \text{GeV}, \quad m_b(1\ \text{GeV}) = 6.89\ \text{GeV}, \quad m_c(m_b) = 1.3\ \text{GeV},$$
$$m_c(2.1\ \text{GeV}) = 1.51\ \text{GeV}, \quad m_s(2.1\ \text{GeV}) = 90\ \text{MeV}, \quad m_s(1\ \text{GeV}) = 119\ \text{MeV},$$
$$m_d(1\ \text{GeV}) = 6.3\ \text{MeV}, \quad m_u(1\ \text{GeV}) = 3.5\ \text{MeV}. \tag{A13}$$

The uncertainty of the strange quark mass is specified as $m_s(2.1\ \text{GeV}) = 90 \pm 20$ MeV.

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