Impact of subleading corrections on hadronic B decays

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We study the subleading corrections originating from the three-parton ($q\overline{q}g$) Fock states of final-state mesons in B decays. The corrections could give significant contributions to decays involving an ω or $\eta^{(\prime)}$ in the final states. Our results indicate the similarity of ωK and $\omega\pi^-$ rates, of order 5×10^{-6}, consistent with the recent measurements. We obtain $a_2(B\to J/\psi K)\approx0.27+0.05i$, in good agreement with data. Without resorting to the unknown singlet annihilation effects, three-parton Fock state contributions can enhance the branching ratios of $K\eta'$ to a level above 50×10^{-6}.

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The rare B decays allow us to access the Kobayashi-Maskawa mixing angles and search for new physics. Much progress in the study of B decays [1–3] has been recently made in QCD-based approaches. In the perturbative QCD (PQCD) framework, the importance of the weak annihilation effects in $B\to K\pi$ decays was first emphasized in Ref. [3], where the annihilation contributions are almost pure imaginary and therefore could lead to CP asymmetry predictions different from the QCD factorization (QCDF) results [1]. Nevertheless, the QCDF study showed that the annihilation effects may play only a minor role in the enhancement of $\pi\pi,\pi K$ branching ratios (BRs) [1]. A recent QCDF fit to $K\pi,\pi\pi$ rates [2] indicated that, even if the annihilation contribution is neglected, one can still get quite good fitting results provided that the strange quark mass is of order 80 MeV.

The annihilation effects might be much more important for VP modes, where P and V denote pseudoscalar and vector mesons, respectively. It has been pointed out that, in the absence of annihilation effects, the ϕK BRs are $\approx4\times10^{-6}$ [4], which is too small compared to the data $\sim8\times10^{-6}$ [5,6]. Recently the Belle experiment observed a large ωK^- rate $(6.7^{+1.3}_{-1.2}\pm0.6)\times10^{-6}$, and $\omega K^-/\omega\pi^-\sim1$ [5]. Sizable ωK results are also reported in new BaBar measurements [6] with $\omega K^{-,0}\sim\omega\pi^-\sim5\times10^{-6}$. It is hard to understand the large strength of ωK rates from the theoretical point of view. The ratio $\omega\overline{K}^0/\omega\pi^-$ reads

$$\omega\overline{K}^0/\omega\pi^-\approx|V_{cb}/V_{ub}|^2(f_K/f_\pi)^2 \times\left|\frac{a_4-a_6r_\chi^K+2r_2(a_3+a_5+a_9/4)+f_Bf_Kb_3}{a_1+r_1a_2}\right|^2, \quad (1)$$

where $r_1=f_\omega F_1^{B\pi}/f_\pi A_0^{B\omega}$, $r_2=(F_1^{BK}f_\omega)/(A_0^{B\omega}f_K)$, $r_\chi^K=2m_K^2/[m_b(m_s+m_u)]$ is the chirally enhanced factor with $m_{s,u}$ being the current quark masses, and $b_3\equiv b_3(K,\omega)$ is the annihilation contribution defined in [7]. The $\omega\pi^-$ rate depends weakly on the annihilation effects. Without annihilation, since a_4 and $a_6r_\chi^K$ terms in the $\omega\overline{K}^0$ amplitude have opposite signs, the ratio $\omega\overline{K}^0/\omega\pi^-$ should be very small. A possibility to explain the data is that the annihilation effects may give the dominant contribution to ωK modes as shown in the QCDF fit [8] for $B\to PP$ together with some $B\to VP$ modes. [However, including the contributions from annihilation effects, the PQCD results read $\mathrm{Br}(\overline{B}^0\to\omega\overline{K}^0)\lesssim2\times10^{-6}$ [9].] This result hints that, to account for the large $\omega\overline{K}^0$ rate, the annihilation contributions to the BRs of all $B\to KV$ modes should be over 80%. If this is true, it should be easy to observe, for instance, the simple relation $\rho^+K^{-,0}:\rho^0K^{-,0}:\omega K^{-,0}\approx1:(1/\sqrt{2})^2:(1/\sqrt{2})^2$, the same as their annihilation ratios squared. Nevertheless, if the global fit is extended to all measured $B\to PV$ modes, a small $K\omega$ rate $\sim2\times10^{-6}$ will be obtained [10] and a reliable best fit cannot be reached. The present QCD approach seems unlikely to offer a coherent picture in dealing with $B\to VP$ modes.

In this article we take into account the subleading corrections arising from the three-parton Fock states of final-state mesons, as depicted in Fig. 1, to QCDF decays amplitudes. We find that it could give significant corrections to decays with ω or $\eta^{(\prime)}$ in the final states. A simple rule extended to $B\to PP,VP$ modes is obtained for the effective coefficients a_i^{SL} with the subleading corrections

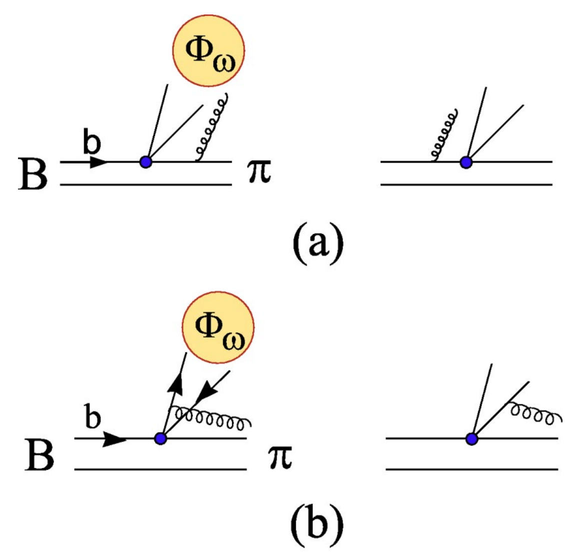

FIG. 1. The contributions of the $q\overline{q}g$ Fock states of the (a) ω and (b) π^- mesons to the $B^-\to\omega\pi^-$ amplitude.

$$a_{2i}^{\mathrm{SL}}=a_{2i}+[1+(-1)^{\delta_{3i}+\delta_{4i}}]c_{2i-1}f_3/2,$$

$$a_{2i-1}^{\mathrm{SL}}=a_{2i-1}+(-1)^{\delta_{3i}+\delta_{4i}}c_{2i}f_3, \tag{2}$$

where $i=1,\ldots,5$, and c_i are the Wilson coefficients defined at the scale $\mu_h=\sqrt{\Lambda_\chi m_B/2}\simeq 1.4$ GeV with Λ_χ the momentum of the emitted gluon as shown in Fig. 1(b) and

$$f_3=\frac{\sqrt{2}}{m_B^2 f_\omega F_1^{B\pi}(m_\omega^2)}\langle\omega\pi^-|O_1|B^-\rangle_{qqg}=0.12 \tag{3}$$

in the SU(3) limit. Here $O_1=\bar{s}\gamma^\mu(1-\gamma_5)u\bar{u}\gamma_\mu(1-\gamma_5)b$, and $\bar{\alpha}_g$ is the averaged fraction of the π^- momentum carried by the gluon. For the ωK amplitudes, the term a_3+a_5, which is originally negligible, is replaced by $a_3+a_5+(c_4-c_6)f_3$, and the latter gives a significantly constructive contribution to the rates. It can thus help in understanding the reason for the similarity between $K\omega$ and $\pi^-\omega$. On the other hand, the subleading corrections can contribute significantly to the processes with $\eta^{(\prime)}$ in the final states, for which the term a_3-a_5 always appears in the decay amplitudes and becomes $a_3-a_5+(c_4+c_6)f_3$ after taking into account the corrections. We also get $a_2^{\mathrm{SL}}(J/\psi K)\approx 0.27+0.05i$ which is well consistent with the data. The result resolves the long-standing sign ambiguity of $\mathrm{Re}(a_2)$.

Let us study the subleading corrections originating from the three-parton Fock states of final-state mesons. Taking the $\omega\pi^-$ mode as an illustration, there are two different types of diagrams shown in Fig. 1. In the following calculation, we adopt the conventions $D_\alpha=\partial_\alpha+ig_sT^aA_\alpha^a$, $\tilde{G}_{\alpha\beta}=(1/2)\epsilon_{\alpha\beta\mu\nu}G^{\mu\nu}$, $\epsilon^{0123}=-1$, and use the Fock-Schwinger gauge to ensure the gauge-invariant nature of the results,

$$A_\mu(x)=-\int_0^1 dv\, v\, G_{\mu\nu}(vx)x^\nu. \tag{4}$$

For Fig. 1(a) where the contributions come from the three-parton Fock states of the ω, because of the $V-A$ structure of the weak interaction vertex, the relevant three-parton light-cone distribution amplitudes (LCDAs) up to the twist-4 level are given by [11]

$$\langle\omega(p_\omega,\lambda)|\bar{u}(0)\gamma_\mu g_s G_{\alpha\beta}(vx)u(0)|0\rangle$$
$$\cong i\frac{f_\omega m_\omega^2}{\sqrt{2}}\int\mathcal{D}\alpha e^{ip_\omega x v\alpha_g}\Big\{(p_\beta^\omega g_{\alpha\mu}-p_\alpha^\omega g_{\beta\mu})\Phi(\alpha_i)$$
$$+\frac{1}{(p_\omega x)}p_\mu^\omega(p_\beta^\omega x_\alpha-p_\alpha^\omega x_\beta)[\Psi(\alpha_i)-\Phi(\alpha_i)]\Big\}, \tag{5}$$

where $\mathcal{D}\alpha=d\alpha_{\bar{u}}d\alpha_u d\alpha_g\delta(1-\alpha_{\bar{u}}-\alpha_u-\alpha_g)$, with $\alpha_{\bar{u}},\alpha_u,\alpha_g$ being the fractions of the ω momentum carried by the $\bar{u}$ quark, u quark, and gluon, respectively. Here Φ and Ψ are the twist-4 LCDAs. Note that all the components of the coordinate x should be taken into account in the calculation before the collinear approximation is applied. The exponential in Eq. (5) before the collinear approximation is actually $e^{ik_g\cdot xv}$, where k_g is the *gluon*'s momentum, and the resultant calculation can be easily performed in the momentum space by substituting $x_\alpha\to-(i/v)(\partial/\partial k_g^\alpha)$. The result of Fig. 1(a) is found to be

$$\langle\omega\pi^-|O_1|B^-\rangle_{\mathrm{Fig.\,1}(a)}$$
$$=f_\omega\frac{4\sqrt{2}m_\omega^2}{3m_B^2}\langle\pi^-|\bar{d}\not{p}_\omega(1-\gamma_5)b|B^-\rangle$$
$$\times\int\mathcal{D}\alpha\frac{2\Phi(\alpha_i)-\Psi(\alpha_i)}{\alpha_g}. \tag{6}$$

Due to G parity, Φ and Ψ are antisymmetric on interchanging $\alpha_{\bar{u}}$ and α_u for the ω, so that Eq. (6) vanishes.

In Fig. 1(b), we consider the emitted gluon which becomes a parton of the pion. We first take $G_{\mu\nu}(vx)\simeq G_{\mu\nu}(0)e^{ivk_g^\pi\cdot x}$ and then adopt the collinear approximation $k_g^\pi=\bar{\alpha}_g p_\pi$ in the final stage of the calculation, where $\bar{\alpha}_g$ is the averaged fraction of the pion's momentum carried by the gluon. The calculation is straightforward and leads to

$$\langle\omega\pi^-|O_1|B^-\rangle_{\mathrm{Fig.\,1}(b)}$$
$$=\frac{f_\omega m_\omega}{4\sqrt{2}N_c}\int_0^1 dv\int_0^1\phi_\omega(u)$$
$$\times\langle\pi^-|\bar{d}\gamma_\mu(1-\gamma_5)g_s\tilde{G}_{\nu\beta}b|B^-\rangle\frac{i\partial}{\partial k_{g\,\beta}^\pi}$$
$$\times\left\{\mathrm{Tr}\left[\not{\epsilon}_\omega^*\left(\frac{\gamma^\nu(v\not{k}_g^\pi+u\not{p}_\omega)\gamma^\mu}{(vk_g^\pi+up_\omega)^2}-\frac{\gamma^\mu(v\not{k}_g^\pi+\bar{u}\not{p}_\omega)\gamma^\nu}{(vk_g^\pi+\bar{u}p_\omega)^2}\right)\right]\right\}$$
$$\cong-\frac{2\sqrt{2}f_\omega}{\bar{\alpha}_g m_B^2}p_\omega^\alpha\langle\pi^-|\bar{d}\gamma^\mu\gamma_5 g_s\tilde{G}_{\alpha\mu}b|B^-\rangle, \tag{7}$$

where the ω mesons's asymptotic leading-twist distribution amplitude $\phi_\omega(u)=6u\bar{u}$ has been taken and $\bar{u}=1-u$. We have two unknown parameters $p_\omega^\alpha\langle\pi^-|\bar{d}\gamma^\mu\gamma_5 g_s\tilde{G}_{\alpha\mu}b|B^-\rangle$ and $\bar{\alpha}_g$ needed to be determined. First, let us evaluate $p_\omega^\alpha\langle\pi^-|\bar{d}\gamma^\mu\gamma_5 g_s\tilde{G}_{\alpha\mu}b|B^-\rangle$. The matrix element can be calculated by considering the correlation function

$$\Pi_\alpha(p,p+q)=i\int d^4x e^{ipx}\langle\pi^-(q)|T[j_{3p}(x)j_B(0)]|0\rangle$$
$$=\frac{m_B^2 f_B}{m_b}\frac{1}{m_B^2-(p+q)^2}$$
$$\times\langle\pi^-|\bar{d}\gamma^\mu\gamma_5 g_s\tilde{G}_{\alpha\mu}b|B^-(q+p)\rangle+\cdots, \tag{8}$$

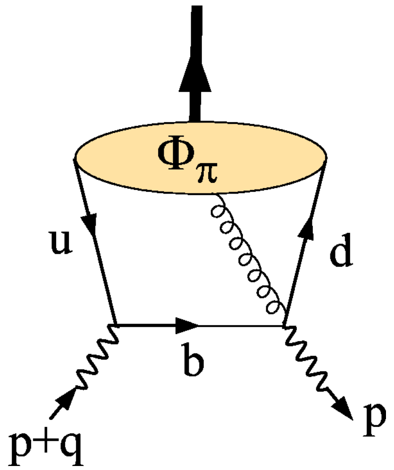

FIG. 2. The diagrammatic illustration to the correlation function, Eq. (8).

where $j_{3p}=\bar{d}g_s\tilde{G}_{\alpha\mu}\gamma^\mu\gamma_5 b, j_B=\bar{b}i\gamma_5 u$, the ellipsis denotes contributions from the higher resonance states, which can couple to the current j_B, and the transition matrix element can be parametrized as

$$\langle\pi^-(q)|\bar{d}\gamma^\mu\gamma_5 g_s\tilde{G}_{\alpha\mu}b|B^-(p+q)\rangle = p_\alpha f_-(p^2)+(p_\alpha+2q_\alpha)f_+(p^2). \quad (9)$$

In the deep Euclidean region of $(p+q)^2$, as depicted in Fig. 2 the correlation function can be perturbatively calculated in QCD and expressed in terms of three-parton LCDAs of the pion,

$$\Pi_\alpha^{\rm QCD}=q_\alpha\int_0^1 \frac{du}{m_b^2-(p+uq)^2}\int_0^u d\alpha_g \times[-2(p\cdot q)f_{3\pi}\phi_{3\pi}+f_\pi m_b(\tilde{\phi}_\parallel-2\tilde{\phi}_\perp)], \quad (10)$$

where $u=\alpha_d+\alpha_g$, and the three-parton pion LCDAs are defined by [12,13]

$$\langle\pi(q)|\bar{d}(x)g_s G_{\mu\nu}(vx)\sigma_{\alpha\beta}\gamma_5 u(0)|0\rangle = if_{3\pi}[q_\beta(q_\mu g_{\nu\alpha}-q_\nu g_{\mu\alpha}) - q_\alpha(q_\mu g_{\nu\beta}-q_\nu g_{\mu\beta})]\int \mathcal{D}\alpha\,\phi_{3\pi}e^{iqx(\alpha_d+v\alpha_g)}, \quad (11)$$

$$\langle\pi(q)|\bar{d}(x)\gamma_\mu g_s\tilde{G}_{\alpha\beta}(vx)u(0)|0\rangle = if_\pi(q_\alpha g_{\beta\mu}-q_\beta g_{\alpha\mu})\int \mathcal{D}\alpha\,\tilde{\phi}_\perp e^{iqx(\alpha_d+v\alpha_g)} - if_\pi\frac{q_\mu}{qx}(q_\alpha x_\beta-q_\beta x_\alpha)\int \mathcal{D}\alpha(\tilde{\phi}_\parallel+\tilde{\phi}_\perp)e^{iqx(\alpha_d+v\alpha_g)}. \quad (12)$$

Here $\phi_{3\pi}$ is a twist-3 DA, and $\tilde{\phi}_\perp$ and $\tilde{\phi}_\parallel$ are all of twist 4,

$$\phi_{3\pi}(\alpha_i)=360\alpha_d\alpha_{\bar{u}}\alpha_g^2\Big[1+\omega_{1,0}\frac{1}{2}(7\alpha_g-3) + \omega_{2,0}(2-4\alpha_d\alpha_{\bar{u}}-8\alpha_g+8\alpha_g^2) + \omega_{1,1}(3\alpha_d\alpha_{\bar{u}}-2\alpha_g+3\alpha_g^2)\Big],$$

$$\tilde{\phi}_\perp(\alpha_i)=30\delta^2\alpha_g^2(1-\alpha_g)\left[\frac{1}{3}+2\varepsilon(1-2\alpha_g)\right],$$

$$\tilde{\phi}_\parallel(\alpha_i)=-120\delta^2\alpha_d\alpha_{\bar{u}}\alpha_g\left[\frac{1}{3}+\varepsilon(1-3\alpha_g)\right]. \quad (13)$$

Since the quark's momentum after emitting the gluon is roughly of order $m_B/2$ and the emitted gluon's momentum is $\Lambda_\chi\sim\bar{\alpha}_g p_\pi$ ($\bar{\alpha}_g$ will be discussed below), we set the scale for the separation of the perturbative and nonperturbative parts at $\mu_h=\sqrt{\Lambda_\chi m_B/2}\simeq 1.4$ GeV. The corresponding parameters at the scale μ_h read $f_{3\pi}=0.0032\ {\rm GeV}^2$, $\omega_{1,0}=-2.63$, $\omega_{2,0}=9.62$, $\omega_{1,1}=-1.05$, $\delta^2=0.19\ {\rm GeV}^2$, $\varepsilon=0.45$ [13]. To calculate $f_\pm$, the contributions of higher resonances in Eq. (8) are approximated by

$$\frac{1}{\pi}\int_{s_0}^\infty \frac{{\rm Im}\Pi_\alpha^{\rm QCD}}{s-(p+q)^2}ds, \quad (14)$$

where s_0 is the threshold of higher resonances. Equating Eqs. (8) and (10) and making the Borel transformation $\mathcal{B}[m_B^2-(p+q)^2]^{-1}=\exp(-m_B^2/M^2)$, we obtain the light-cone sum rule

$$f_-(p^2)=\frac{m_b}{2m_B^2 f_B}\int_0^1 du\int_\Delta^u d\alpha_g e^{([m_B^2/M^2-(m_b^2-\bar{u}p^2)/uM^2])} \times\left[f_{3\pi}\frac{m_b^2-p^2}{u^2}\phi_{3\pi}-f_\pi\frac{m_b}{u}(\tilde{\phi}_\parallel-2\tilde{\phi}_\perp)\right], \quad (15)$$

and $f_+(p^2)=-f_-(p^2)$, where $\Delta=u-(m_b^2-p^2)/(s_0-p^2)$. Using the above parameters for LCDAs, $m_b=4.7\pm0.1$ GeV, and $f_B=180$ MeV, we obtain the stable $f_\pm$ prediction by adopting $s_0\simeq 37\ {\rm GeV}^2$ and $M^2\approx[9-20]\ {\rm GeV}^2$. The result is depicted in Fig. 3(a). The resulting value is $f_\mp(m_\omega^2)=\pm(0.057\pm0.005)\ {\rm GeV}^2$, where the uncertainty comes from the sum rule analysis. We then get

$$p_\omega^\alpha\langle\pi^-|\bar{d}\gamma^\mu\gamma_5 g_s\tilde{G}_{\alpha\mu}b|B^-\rangle=-f_-(m_\omega^2)\times(m_B^2-m_\pi^2-m_\omega^2) \simeq -1.6\ {\rm GeV}^4. \quad (16)$$

Next, we determine the value of $\bar{\alpha}_g$. For illustration, we plot in Fig. 3(b) the amplitude A_{f_-} of the $f_-(m_\omega^2)$ sum rule versus α_g and $u(=\alpha_d+\alpha_g)$ by adopting $M^2=10\ {\rm GeV}^2$, where A_{f_-} satisfies $f_-=\int_0^1 du\int_0^u d\alpha_g A_{f_-}$, i.e., the volume in the plot is equal to $f_-(m_\omega^2)$. The resultant form factor is dominated by the region where $u\gtrsim 60\%, \alpha_g\lesssim 30\%$. The averaged fraction of the pion momentum carried by the gluon is then estimated to be $\bar{\alpha}_g=(\int_0^1 du\int_0^u d\alpha_g\alpha_g A_{f_-})/f_-\simeq 0.23$.

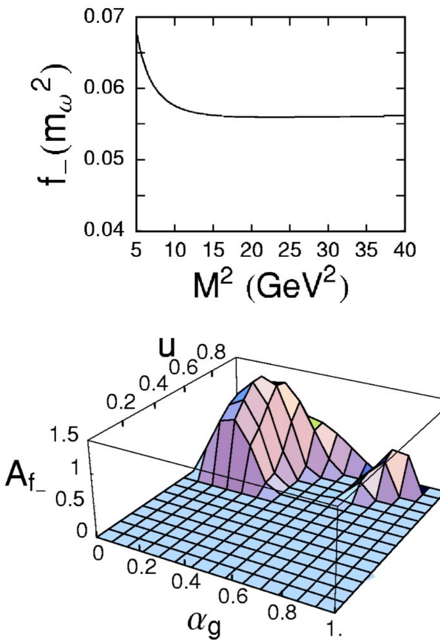

FIG. 3. (a) Form factor $f_-(m_\omega^2)$ plotted as a function of the Borel mass squared M^2. (b) $f_-(m_\omega^2)=\int_0^1 du \int_0^u d\alpha_g A_{f_-}$ with $M^2=10$ GeV^2. The volume in the plot is equal to $f_-(m_\omega^2)$. Here $u=\alpha_d+\alpha_g$.

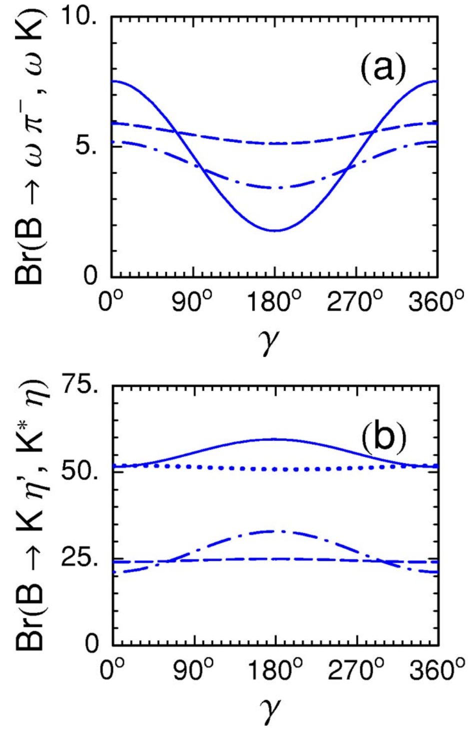

FIG. 4. (a) Dashed, solid, and dot-dashed lines for $\overline{B}\to\omega\pi^-,\omega K^-$ and $\omega\overline{K}^0$; (b) solid, dotted, dot-dashed, and dashed lines for $K^-\eta',\overline{K}^0\eta',K^{*-}\eta$, and $\overline{K}^{*0}\eta$. BRs are in units of 10^{-6}.

We therefore obtain $\langle\omega\pi^-|O_1|B^-\rangle_{\text{Fig. 3}(b)}\simeq 0.13$ GeV^{-3} and $f_3=0.12$ which gives the correction to a_i as defined in Eq. (2). We list a_i without and with the subleading corrections in Table I, where the approximation $-\sqrt{2}ip_\pi^\alpha\langle\omega|\overline{u}\gamma^\mu g_s\widetilde{G}_{\alpha\mu}b|B^-\rangle/A_0^{B\omega}\simeq p_\omega^\alpha\langle\pi^-|\overline{d}\gamma^\mu\gamma_5 g_s\widetilde{G}_{\alpha\mu}b\times|B^-\rangle/F_1^{B\pi}$ has been made. Note that a_6 and a_8 do not receive subleading corrections.

In the following analysis, the LC sum rule form factors and $m_s=80$ MeV are used. We will instead use a smaller $A_0^{B\rho}=0.28$ which is preferred by the $\omega\pi^-$ data. The $\omega\pi^-$ mode is ideal for extracting $A_0^{B\rho}$ since its rate is insensitive to annihilation effects. We find that the spectator parameter $X_H=[\ln(m_B/\Lambda_h)](1+\rho_H e^{i\phi_H})$ is consistent with zero in the analysis. The reason is that, since the spectator interaction with a gluon exchange between the emitted meson and the recoiled pseudoscalar meson of twist-3 LCDA Φ_σ is endpoint divergent in the collinear expansion, the vertex of the gluon and spectator quark should be considered inside the pion wave function, i.e., for this situation the pion itself is in a three-parton Fock state. Annihilation effects have been emphasized in ϕK studies [4]. We adopt the annihilation parameters $\rho_A\simeq 0.9,\phi_A\simeq 0$ which give $\text{Br}(B^-\to\phi K^-)\simeq 8.5\times 10^{-6}$, consistent with the current data. Here the annihilation parameter of VP modes is defined as $X_A^{VP}=[\ln(m_B/\Lambda_h)](1+\rho_A e^{i\phi_A})$ [1] and its imaginary part is neglected since the BRs are insensitive to it. In Fig. 4(a) we plot the BRs of $\omega\pi^-$ and ωK modes versus γ ($\equiv\arg V_{ub}^*$). The results for $\gamma\approx(60\text{–}120)°$ are in good agreement with data. At $\gamma=90°$, the plot gives $\omega\pi^-,\omega K^-,\omega\overline{K}^0$ to be 5.5, 4.5, 4.3, respectively, in units of 10^{-6}. Without the contributions from three-parton Fock states of mesons, $\phi K^-,\omega\pi^-,\omega K^-,\omega\overline{K}^0$ will become 11, 3.9, 3.1, 2.9 (in units of 10^{-6}). The three-parton Fock state effects give constructive contributions to $\omega\pi,\omega K$ modes, but a destructive one to the ϕK mode.

The corrections from three-parton Fock states of the kaon also give a definite answer to the long-standing problem for $a_2(J/\psi K)$. In the earlier study, to account for the experimental value $|a_2|$, the parameter ρ_H has to be $\gtrsim 1.5$ [14]. As emphasized in passing, without fine-tuning ρ_H, we calculate the amplitudes from the three-parton Fock states of the kaon. With the same procedure as shown above, we obtain $f_\mp(m_{J/\psi}^2)\simeq\pm 0.08,f_3\simeq 0.14$, and thus $a_2^{\text{SL}}=a_2^{t2}+c_1(\mu_h)f_3\simeq 0.10+0.05i+c_1(\mu_h)0.14=0.27+0.05i$, where a_2^{t2} is de-

TABLE I. Values for a_i for charmless B decay processes without (first row) and with (second row) three-parton Fock state contributions of final-state mesons, where a_3–a_{10} are in units of 10^{-4} and the annihilation effects are not included.

a_1	a_2	a_3	a_4	a_5	a_6	a_7	a_8	a_9	a_{10}
$1.02+0.014i$	0.10-$0.08i$	$26+26i$	-328-$91i$	1.2-$30i$	-487-$72i$	$0.7+0.3i$	$4.5+0.6i$	-89-$0.1i$	$-5.9+7i$
$0.974+0.014i$	0.25-$0.08i$	$-55+26i$	-291-$91i$	112-$30i$	-487-$72i$	$-0.3+0.3i$	$4.5+0.6i$	-88-$0.1i$	$-18+7i$

termined up to the twist-2 order and the SU(3) approximation for f_3 has been made. The result for $a_2^{\rm SL}$ is well consistent with that extracted from data. This solves the long-standing sign ambiguity of $a_2(J/\psi K)$ which turns out to be positive for its real part. Note that if $\mathrm{Re}(a_2)$ were negative, f_3 would have to be ~ -0.3, which in turn would lead to $\phi K \sim 20\times 10^{-6}$ and $\omega\pi^-, \omega K \sim 1\times 10^{-6}$.

The subleading corrections could give significant contributions to the decays with $\eta^{(\prime)}$ in the final states because these decay amplitudes always contain the singlet factor $a_3 - a_5$. We plot the BRs of $K\eta', K^*\eta$ modes versus γ in Fig. 4(b), where $X_A^{PP} \approx 0$, the annihilation parameter for *PP* modes, has been used as it could give good fit results for $K\pi, \pi\pi$ rates [2]. We do not consider the singlet annihilation correction [15] in $K\eta'$ modes because it is still hard to determine at present. With (without) the subleading corrections, we see that $K^-\eta' \gtrsim \overline{K}^0\eta' \approx 55(35)$, and $K^{*-}\eta \gtrsim \overline{K}^{*0}\eta \approx 24(20)$, in units of 10^{-6}. The corrections give 70% and 25% enhancements to the $K\eta'$ and $K^*\eta$ rates, respectively. Note that with (without) the corrections, $K^*\eta' \lesssim 1\times 10^{-6}$ ($\gtrsim 4\times 10^{-6}$) and $K\eta \lesssim 1\times 10^{-6}$ ($\gtrsim 1\times 10^{-6}$). In a PQCD calculation [16], it seems that $\overline{K}^0\eta'$ ($=41\times 10^{-6}$) was underestimated while $\overline{K}^0\eta$ ($=7\times 10^{-6}$) was overestimated compared to the data [5,6]. Within the QCDF framework, by only considering the two-parton LCDAs of mesons, Beneke and Neubert [15] obtained $K^*\eta \sim 13\times 10^{-6}$, just half of the experimental value, but with a huge error.

For further comparison with other calculations, we plot the $K\rho, K^*\pi, \rho^-\eta^{(\prime)}$ modes in Fig. 5. The subleading contributions to these BRs are $\lesssim 15\%$. At $\gamma = 90°$, we have $\overline{K}^0\rho^0, K^-\rho^+, \overline{K}^0\rho^-, K^-\rho^0 = 6,7,7,2(\times 10^{-6})$, and $K^{*-}\pi^+, \overline{K}^{*0}\pi^0, \rho^-\eta, \rho^-\eta' = 9,2,6,4(\times 10^{-6})$. In [7], to fit $\phi K, \omega K$ rates, the annihilation effects dominate the decay amplitudes and the form factors $F_1^{BK}, A_0^{B\pi}$ are rather small, such that it will lead to $\rho^+K^{-,0}:\rho^0K^{-,0}:\omega K^{-,0} \approx 1:(1/\sqrt{2})^2:(1/\sqrt{2})^2$, while the PQCD results for the rates are $\overline{K}^0\rho^0, K^-\rho^+, \overline{K}^0\rho^-, K^-\rho^0 = 2.5,5.4,3.0,2.2(\times 10^{-6})$ [9].

In conclusion, we have calculated the contributions arising from three-parton Fock states of mesons in *B* decays. We find that the contributions could give significant corrections to decays with an ω or $\eta^{(\prime)}$ in the final states. Our main results for $\gamma \approx (60\text{–}110)°$ are $\omega\pi^-, \omega K^-, \omega\overline{K}^0 \approx 6.0,(6\text{–}5),5.1$, respectively, in units of 10^{-6}, while the previous QCDF global fit to *VP* modes and the pQCD results gave smaller ωK BRs of order $\lesssim 3\times 10^{-6}$ [9,10]. We predict that $\overline{K}^0\rho^0 \sim \omega K$ and $\overline{K}^0\rho^0/K^-\rho^0 \approx 3$. Including the corrections, we obtain $a_2(J/\psi K) \approx 0.27 + 0.05i$, which is well consistent with the data. The sign of $\mathrm{Re}(a_2)$ turns out to be positive. Without resorting to the unknown singlet annihilation effects, three-parton Fock state contributions can enhance $K\eta'$ to a level above 5×10^{-5}.

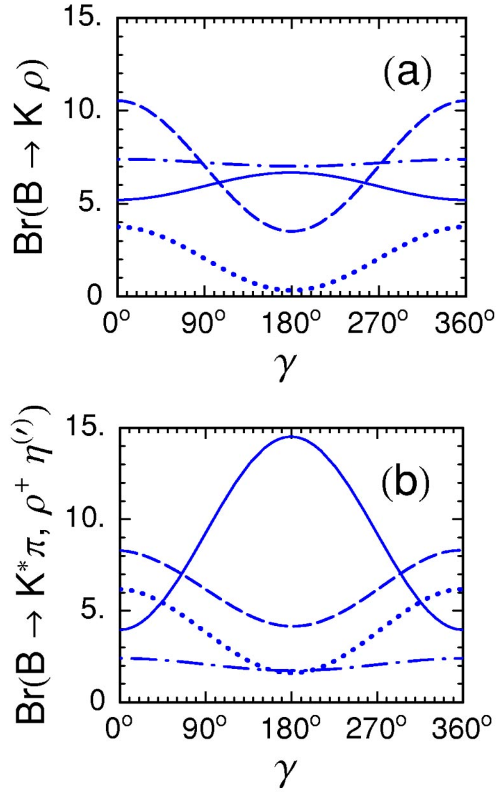

FIG. 5. (a) Solid, dashed, dot-dashed, and dotted lines for $\overline{B} \to \overline{K}^0\rho^0, K^-\rho^+, \overline{K}^0\rho^-$, and $K^-\rho^0$; (b) solid, dot-dashed, dashed, and dotted lines for $\overline{B} \to K^{*-}\pi^+, \overline{K}^{*0}\pi^0, \rho^-\eta$ and $\rho^-\eta'$. BRs are in units of 10^{-6}.

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