Importance of Higher Twist Effects to Understand Charmed
Color-Suppressed $B$ Decays

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Abstract

Working within the framework of the QCD light cone sum rules (LCSR), we compute and discuss the nonfactorizable higher twist effect in $\bar{B}^0 \rightarrow D^{*0} \pi^0$ to make an all-around examination of its role in the charmed color-suppressed $B$ decays. Analogously to the case of $\bar{B}^0 \rightarrow D^0 \pi^0$, such effect turns out to be of the same strong phase as the factorizable amplitude, and modifies constructively the magnitude by order $(40 - 90)\%$ so that the effective coefficient $a^f_2 = C_1 + C_2/3$ receives a positive correction comparable numerically with it. Nonleading as the soft effect in question is, our findings for it, along with the previous LCSR analyses of $\bar{B}^0 \rightarrow D^0 \pi^0$, are suggestive of the dominance of soft exchanges in these charmed color suppressed $B$ decays. Also, the emphases are put on importance of understanding intensively various related higher twist and transverse momentum effects to interpret the data on $B \rightarrow D^{0(*)} \left(\pi^0, \eta, \eta'\right), J/\psi K^{(*)}$. 

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The formulation of the QCD factorization (QCDF) \[\text{[1]}\] and soft-collinear effective theory (SCET) \[\text{[2]}\] has greatly renewed and deepened our understanding of QCD dynamics in the two-body hadronic decays of B-mesons, and as such considerably enhanced our confidence in precisely estimating the decay amplitudes and CP-asymmetry effects. Especially, ones have essentially changed minds in understanding of the charmed color-suppressed B decays and realized that the nonfactorizable soft exchange contributions dominate over the hard-exchange ones. For all these progresses, however, being unable to deal with non-perturbative QCD dynamics from the first principle makes yet it difficult to theoretically interpret the experimental observations of the color suppressed process $\bar{B}^0 \to D^{0(\ast)}\pi^0$ \[\text{[3]}\]. One attempts to make sense of them by a global analysis of the data on $B \to D^{(\ast)}\pi$ decays, either in a model-independent manner \[\text{[4]}\] \[\text{[5]}\] \[\text{[6]}\], or in the framework of SCET \[\text{[7]}\] and Perturbative QCD \[\text{[8]}\]. The leading soft effects in such decays are supposed to include final-state rescattering processes and nonfactorizable contributions induced by the color-octet operators. Resorting to various phenomenological approaches, such as one-particle exchange model \[\text{[9]}\] and Regge theory \[\text{[10]}\], one models the soft contributions from the former. QCD sum rule approach is initially used \[\text{[11]}\] and QCD light-cone sum rule (LRSR) method \[\text{[12]}\] is subsequently applied \[\text{[13]}\] to obtain an estimate of the latter. Focusing on the $\bar{B}^0 \to D^{0}\pi^0$ case, the author of \[\text{[11]}\] finds that the soft exchanges between the color-octet quark pair $c\bar{q}$ and the pion, which concern higher-twist components of the pion, provide a negative correction at a level of 70% of the factorizable amplitude, while the LCSR approach predicts a positive numerical effect at $(50 - 110)\%$ order \[\text{[14]}\]. As compared to the $B \to J/\Psi K$ case where the effect in question contributes to the magnitude by $(30 - 70)\%$ \[\text{[15]}\], the higher-twist effect in $\bar{B}^0 \to D^{0}\pi^0$ is, as it were, more important, as expected. Also, the similar calculations are carried out for $B \to \pi\pi$ \[\text{[13]}\] \[\text{[16]}\] and $B \to K\pi$ \[\text{[17]}\], but the resulting numerical impacts are found to be far less than those in both $\bar{B}^0 \to D^{0}\pi^0$ and $B \to J/\Psi K$ cases, a result in accordance with the argument for QCD factorization.

Whereas the nonfactorizable higher-twist effect, as estimated, plays a role comparable numerically with the factorizable amplitudes of the color-suppressed $\bar{B}^0 \to D^{0}\pi^0$ and $B \to J/\Psi K$ decays, in the letter we would like to examine how it affect $\bar{B}^0 \to D^{0\ast}\pi^0$, to get an all-around understanding of its role in the charmed color suppressed $B$ decays.

We start with looking briefly back at the dynamical features of the color suppressed $\bar{B}^0 \to D^{0(\ast)}\pi^0$. In the Fierz transformed form, the relevant effective weak Hamiltonian reads \[\text{[18]}\]

$$\mathcal{H}_W \equiv \frac{G_F}{\sqrt{2}} V_{cb}V_{ud}^\ast \left[(C_1(\mu) + \frac{1}{3}C_2(\mu))O_1(\mu) + 2C_2(\mu)\tilde{O}_1(\mu)\right], \quad (1)$$

where $C_{1,2}$ are the Wilson coefficients, $V_{ij}$ the CKM matrix elements, and

$$O_1 = (\bar{c}\Gamma^\mu u)(\bar{d}\Gamma_\mu b), \quad \tilde{O}_1 = (\bar{c}\frac{\lambda_\alpha}{2}\Gamma_\mu u)(\bar{d}\frac{\lambda_\alpha}{2}\Gamma^\mu b), \quad (2)$$

the color-single and -octet quark operators, respectively, with $\Gamma_\mu = \gamma_\mu (1 - \gamma_5)$ and $\lambda_\alpha$ being the color $SU(3)$ matrices. It is clear that the transitions induced by $\tilde{O}_1$ come from
gluon exchanges and thus the corresponding matrix elements are non-factorizable. Since the ejected \( \bar{D}^{0(*)} \) in the decays is a heavy meson inside which the two quarks are asymmetric in momentum distribution, the non-factorizable vertex diagrams prove themselves to be QCD infrared divergent so that QCD factorization breaks down \[19\]. This hints at the dominance of soft dynamics. If we introduce an effective coefficient \( a_2 \) such that the decay amplitude of \( \bar{B}^0 \to D^{*0}\pi^0 \) is made be of the following parametrization:

\[
\mathcal{A}(\bar{B}^0(p + q) \to D^{*0}(\epsilon, p)\pi^0(q)) = -G_F V_{cb} V_{ud}^* \epsilon \cdot q m_{D^*} f_{D^*} F_1^{B \to \pi}(m_{D^*}) a_2, \tag{3}
\]

with \( \epsilon, m_{D^*} \) and \( f_{D^*} \) being the polarization vector, mass and decay constant of \( D^* \) meson, respectively, and \( F_1^{B \to \pi}(m_{D^*}) \) the \( B \to \pi \) form factor defined as

\[
\langle \pi(q) \mid \bar{u}\gamma_\mu b \mid B(p + q) \rangle = (2q + p)_\mu F_1^{B \to \pi}(p^2) + p_\mu \frac{m_B^2 - m_\pi^2}{p^2} (F_0^{B \to \pi}(p^2) - F_1^{B \to \pi}(p^2)), \tag{4}
\]

\( a_2 \) is actually supposed to absorb all the nonfactorizable contributions including higher twist component we are to calculate.

Let’s go over to LCSR calculation of the nonfactorizable higher-twist effect involved in \( \bar{B}^0 \to D^{*0}\pi^0 \). With this end in view, we opt for the following correlator:

\[
\Pi_\mu(p, q, k) = i^2 \int \int d^4x d^4ye^{-i(p+q)x+i(p-k)k} \times \langle \pi^0(q)\mid T\{\bar{u}(y)\gamma_\mu c(y), \bar{O}_1(0), m_{b}(x)i\gamma_5d(x)\}\rangle|0\rangle, \tag{5}
\]

where the current operators \( \bar{u}\gamma_\mu c \) and \( \bar{b}\gamma_5d \) are used to interpolate \( D^{*0} \) and \( B^0 \) meson fields respectively, and an unphysical momentum \( k \) is introduced to make \( B \) channel free of pollution by \( D^{*}\pi \) resonances, but it must disappear from the physical \( \bar{B}^0 \to D^{*0}\pi^0 \) amplitude in the dispersion integral. Schematically, the higher twist contributions to the correlator can be illustrated by Feynman diagrams in Fig.1. The correlation function \( \Pi_\mu \) is a functions of 6 independent invariants, which, adapting to the present purpose, are to be chosen as \( p^2, q^2, (p + q)^2, (p - k)^2, k^2 \) and \( P^2 = (p + q - k)^2 \). At small light-cone distances \( x^2 \sim y^2 \sim (x - y)^2 \sim 0 \), or equivalently in large and space-like regions of \( P^2, (p + q)^2 \) and \( (p - k)^2 \), the Operator Product Expansion (OPE) is applicable to Eq.\[5\]. Furthermore, not violating the validity of OPE we simply let \( k^2 = 0 \) to the pion.

Inserting on the right hand side of \[5\] a complete set of hadronic states with \( D^{*0} \) quantum numbers, we obtain the phenomenological expression \( \Pi_\mu^{H(D^{*})} \),

\[
\Pi_\mu^{H(D^{*})}((p-k)^2, (p+q)^2, P^2, p^2) = \frac{m_{D^{*0}} f_{D^{*0}}}{m_{D^*}^2 - (p-k)^2} G_\mu((p+q)^2, P^2, p^2) + \int_{s_0^{H(D^{*})}}^\infty ds \frac{\rho_\mu^{H(D^{*})}(s, (p+q)^2, P^2, p^2)}{s - (p-k)^2}, \tag{6}
\]
where
\[ G_\mu = i \sum_\lambda e^\lambda_\mu \int d^4x \ e^{-i(p+q)x} \langle D^{00}(\epsilon^\lambda, p-k) \pi^0(q)|T\{\bar{O}_1(0), \ m_b\bar{b}(x)i\gamma_5d(x)\}|0 \]
\[ = \sum_\lambda e^\lambda_\mu \Pi((p+q)^2, P^2, P^2, \epsilon^*(\lambda) \cdot Q_\iota) \]
\( (7) \)

is a two point function with \( Q_\iota = p, q \) and \( k, \ \rho^{H(D^*)}_\mu \) and \( s_0^{H(D^*)} \) are respectively the hadronic spectral function and threshold mass squared of the higher states in \( D^* \) channel. The integral of \( (6) \), invoking quark-hadron duality, is supposed to coincide with the \( s \geq s_0^{D^*} \) part of the QCD spectral representation of \( \Pi_\mu \)
\[ \frac{1}{\pi} \int_{s_0^{D^*}}^\infty ds \ \text{Im} \Pi_{\mu}^{(QCD)}(s, (p+q)^2, P^2, P^2) \frac{s-(p-k)^2}{s} \]
\( (8) \)

with \( s_0^{D^*} \) being the QCD effective threshold in \( D^* \) channel, and \( \frac{1}{\pi} \text{Im} \Pi_{\mu}^{(QCD)} \) the related QCD spectral density. Then we equate \( (6) \) with the QCD answer for \( (5) \) and make the Borel transformation \( (p-k)^2 \to T \), yielding
\[ G_\mu(p+q)^2, P^2, P^2 = \frac{1}{\pi m_{D^*} f_{D^*}} \int_{s_0^{D^*}}^{s_{D^*}} ds \ e^{(m_{D^*}^2-s)/T} \text{Im} \Pi_{\mu}^{(QCD)}(s, (p+q)^2, P^2, P^2). \]
\( (9) \)

To proceed, saturating the two-point function \( G_\mu \) with the intermediate states of \( B^0 \) quantum numbers, we have its hadronic expression in which the matrix element \( \langle D^{00}(\epsilon^\lambda, p-k)\pi^0(q)|\bar{O}_1|B^0(p+q) \rangle \) enters the lowest pole term. When the condition \( (p+q)^2 = (p+q-k)^2 = m_B^2 \) is imposed on the \( B^0(p+q) \to D^{00}(\epsilon^\lambda, p-k)\pi^0(q) \) transition, the unphysical momentum \( k \) vanishes automatically so that the corresponding matrix element is made physical. On some algebraic manipulations we see easily that the required physical amplitude can be parameterized in term of a function factor \( M^{(p-k)} \) in front of \( (p-k)_\mu \) in the kinematic decomposition of \( \sum_{\epsilon^\lambda} \langle D^{00}(\epsilon^\lambda, p-k)\pi^0(q)|\bar{O}_1|B^0(p+q) \rangle \). This denotes that we need only to calculate the related part \( \Pi_{\mu}^{(p-k)} \) of \( \Pi_{\mu}^{(QCD)} \). Furthermore, to apply quark hadron duality to \( B \) channel we stick down \( \text{Im} \Pi_{\mu}^{(QCD)}(s, (p+q)^2, P^2, P^2) \) in a form of dispersion integral in variable \( (p+q)^2 \), and make the analytic continuation \( P^2 \to m_B^2, \) keeping \( (p+q)^2 \) fixed and letting \( p^2 = m_B^2 \). Then using the standard procedure the desired sum rule can be achieved, if the relevant QCD double spectral density \( \rho_{QCD}(s, s', m_B^2, p^2) \) is known, where \( s' \) is a spectral variable in \( B \) channel.

Now we embark on the computation of \( \Pi_{\mu}^{(p-k)} \) to extract \( \rho_{QCD} \). It is advisable that we work with the fixed-point gauge in which the light cone expansion of quark propagator is taken in a correction term due to the interaction with one background-field gluon. After some lengthy calculations we obtain the twist-3 contribution:
\[ \Pi_{\mu}^{(p-k)} = \frac{m_b f_{3\pi}}{8\sqrt{2}\pi^2} \int_{s_0^{D^*}}^\infty ds \ \text{Im} \Pi_{\mu}^{(QCD)}(s, (p+q)^2, P^2, P^2) \]
\[ \int_0^1 dt \int_0^1 du \int_{m_b^2}^{u} dv \frac{1}{v^2} \frac{1}{s} \]
\[ \frac{ds}{s-(p-k)^2} \int_{\chi(s,t,m_B^2)}^1 du \int_{\chi(s,t,m_B^2)}^1 dv \frac{1}{v^2} \]
\[ \frac{m_b^2-(p+q)^2}{v^2} \]
\[ \]
\[ \times \phi_{3\pi}(1 - u, u - v, v) \left\{ (s - \frac{m_c^2}{t})(1 - t) + 2t[(p + q)^2 - p^2] \right\} \times (2v - \chi(s, t, m_B^2)) \] 

In the above equation, \( f_{3\pi} \) is a nonperturbative quantity defined by the vacuum-pion matrix element of the local operator \( \bar{u}\sigma_{\mu\nu}\gamma_5 G_{\alpha\beta}d \), \( \phi_{3\pi} \) refers to a pionic twist-3 distribution amplitude related to the nonlocal operator \( \bar{d}(0)\sigma_{\mu\nu}\gamma_5 G_{\alpha\beta}(vy)u(x) \) [21], and \( \chi = (s - \frac{m_c^2}{t})/(s - m_B^2) \). The fact that the threshold \( s_0^D \) is far below the squared mass of b quark allows us to make an expansion of (10) in \( \chi \). Then passing over those terms vanishing after Borel transformations we have, up to order \( \mathcal{O}(\chi^3) \),

\[ \rho_{tw3}(s, s', m_B^2, p^2) = \frac{m_b f_{3\pi}}{16\sqrt{2} \pi^2 (s' - p^2)} \int_{m_c^2}^{1} dt \times \left\{ 2 \int_{0}^{u} \frac{dv}{v^2} \phi_{3\pi}(1 - u, u - v, v) \left[ (1 - t) \left( s - \frac{m_c^2}{t} \right) + 2vt(s' - p^2) \right] 
+ 2 \left[ (1 - 3t)(s' - p^2) \int_{0}^{u} \frac{dv}{v^2} \phi_{3\pi}(1 - u, u - v, v) 
+ (t - 1) \left( s - \frac{m_c^2}{t} \right) \left( \frac{1}{v^2} \phi_{3\pi}(1 - u, u - v, v) \right)_{v=0} \right] \chi(s, t, m_B^2) 
+ \left[ (t - 1) \left( s - \frac{m_c^2}{t} \right) \frac{\partial}{\partial v} \left( \frac{1}{v^2} \phi_{3\pi}(1 - u, u - v, v) \right) 
+ 2(2t - 1)(s' - p^2) \frac{1}{v^2} \phi_{3\pi}(1 - u, u - v, v) \right] \chi \right\}. \] 

with \( u = (m_b^2 - p^2)/(s' - p^2) \). Following the same line the twist-4 contribution is attainable. Here we do not give it any more to save some space.

At the end, the required nonfactorizable effect in \( \bar{B}^0 \to D^{*0}\pi^0 \) can be numerically estimated by the following LCSR answer:

\[ \langle D^{*0}(\epsilon, p)\pi^0(q)| \bar{O}_1(0)| \bar{B}^0(p + q) \rangle = -\frac{\epsilon^* \cdot q m_b m_{D^*}}{4\sqrt{2} \pi^2 f_{D^*} f_B m_B^2 (m_B^2 - m_{D^*}^2)} \int_{m_c^2}^{s_0^B} ds \; e^{(m_{D^*}^2 - s)/T} \times \int_{u_0}^{1} du \; e^{(m_B^2 - (m_B^2 - m_{D^*}^2)(1-u))/T} \int_{m_c^2}^{1} dt \left\{ f_{3\pi} \left[ \int_{0}^{u} \frac{dv}{v^2} \phi_{3\pi}(1 - u, u - v, v) 
\times \left( \frac{m_B^2 - m_{D^*}^2}{u} \right) [2vt + (1 - 3t)\chi] + \left( s - \frac{m_c^2}{t} \right)(1 - t) \right] 
+ (t - 1) \left( s - \frac{m_c^2}{t} \right) \left[ \frac{1}{v^2} \phi_{3\pi}(1 - u, u - v, v) \right]_{v=0} \chi \right\}. \]
+ \left( 2(2t - 1)\frac{m_b^2 - m_{D*}^2}{u} \frac{1}{v^2} \phi_{3\pi}(1 - u, u - v, v) \right) \\
+ t - 1 \left( s - \frac{m_c^2}{t} \right) \frac{\partial}{\partial v} \left[ \frac{1}{v^2} \phi_{3\pi}(1 - u, u - v, v) \right]_{v=0} \chi^2 \\
+ m_b f_\pi \int_0^u \frac{dv}{v^2} \left[ 3v \left[ (2t - 1)\phi_{\perp}(1 - u, u - v, v) + \tilde{\phi}_{\perp}(1 - u, u - v, v) \right] \\
- 2 \left[ (2t - 1)\phi_{\perp}(1 - u, u - v, v) + \tilde{\phi}_{\perp}(1 - u, u - v, v) \right] \chi \right] \\
+ \frac{1}{2} \left( \frac{3}{v^2} \left[ (2t - 1)\phi_{\perp}(1 - u, u - v, v) + \tilde{\phi}_{\perp}(1 - u, u - v, v) \right] \\
+ \frac{1}{v} \frac{\partial}{\partial v} \left[ (2t - 1)\phi_{\perp}(1 - u, u - v, v) - \tilde{\phi}_{\perp}(1 - u, u - v, v) \right] \right]_{v=0} \chi^2 \right) \right), \quad (12)

where \( n_B^0 = (m_B^2 - m_{D*}^2)/(s_B^0 - m_{D*}^2) \) with \( s_B^0 \) being the QCD effective threshold in \( B \) channel, \( T' \) indicates the Borel parameter corresponding to the variable \((p + q)^2\), the pionic distribution amplitude \( \phi_{\perp}(\tilde{\phi}_{\perp}) \) is of twist-4 \[21\].

To go on with numerical discussions. We take \[22, 23\] \( m_{D*}^0 = 2.007 \text{ GeV} \), \( m_c = 1.3 \pm 0.1 \text{ GeV} \), \( f_{D*} = 240 \pm 10 \text{ MeV} \), \( s_B^0 = 6 \pm 1 \text{ GeV}^2 \) and \( T = 1 - 2 \text{ GeV}^2 \) for the \( D^* \) channel parameters, and \( m_B = 5.28 \text{ GeV} \), \( m_b = 4.7 \pm 0.1 \text{ GeV} \), \( f_B = 180 \pm 30 \text{ GeV} \), \( s_B^0 = 35 \pm 2 \text{ GeV}^2 \) and \( T' = 8 - 12 \text{ GeV}^2 \) for the \( B \) channel ones. The nonperturbative quantities concerning \( \pi \) meson are chosen as \[23\]: \( f_\pi = 132 \text{ MeV} \), \( f_{3\pi}(\mu_b) = 0.0026 \text{ GeV} \), and

\[
\varphi_{3\pi}(\alpha_i, \mu_b) = 360\alpha_1\alpha_2\alpha_3^2 \left[ 1 + \frac{\omega_{3\pi}}{2}(7\alpha_3 - 3) \right], \\
\varphi_{\perp}(\alpha_i, \mu_b) = 30\delta_3^2(\alpha_1 - \alpha_2)\alpha_3^2 \left[ \frac{1}{3} + 2\epsilon_{3\pi}(1 - 2\alpha_3) \right], \\
\tilde{\varphi}_{\perp}(\alpha_i, \mu_b) = 30\delta_3^2(1 - \alpha_3)\alpha_3^2 \left[ \frac{1}{3} + 2\epsilon_{3\pi}(1 - 2\alpha_3) \right],
\]

with \( \mu_b = \sqrt{m_B^2 - m_b^2} \sim m_b/2 \), a scale characteristic of the typical virtuality of the underlying \( b \) quark, and \( \omega_{3\pi} = -0.28, \delta_3^2 = 0.17 \text{GeV}^2, \epsilon_{3\pi} = 0.36 \). The other parameters are fixed as \( C_1(\mu_b) = -0.26, C_2(\mu_b) = 1.12 \) and \( F_{1B \to \pi}(m_{D*}^2) = 0.30 \[22\].

Using these inputs, we see that with both the Borel parameters varying in the respective intervals, Eq. (12) reveals a good stability in numerical value. From the sum rule \textit{windows}, the resulting impact on the decay amplitude could be derived using

\[
A_S(\bar{B}^0 \to D^{*0}\pi^0) = \sqrt{2}G_F V_{cb} V_{ud}^* C_2(D^{*0}(\epsilon, p)\pi^0(q))|\tilde{O}_1(0)|\bar{B}^0(p + q)).
\]

Nevertheless, for the purpose of comparing with the factorizable amplitude and reducing the uncertainty, it is more advisable to calculate the ratio of \( A_S \) over the factorizable

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amplitude $A_F$. The result is

$$R = A_S(\bar{B}^0 \to D^{*0}\pi^0)/A_F(\bar{B}^0 \to D^{*0}\pi^0).$$

$$= 0.43 - 0.95. \quad (15)$$

The range quoted here is due to the variations of the various inputs and is obtained by adding linearly up the corresponding uncertainties. The soft contribution is shown to be of the same phase as the factorizable one, indicating that it does not provide the decay amplitude with any strong phase, and with a certain uncertainty such effect modifies the magnitude of the factorizable amplitude at a comparable or even identical level.

When a comparison is drawn with the situation of $\bar{B}^0 \to D^0\pi^0$ in which the ratio is estimated at between 0.54 and 1.15 [14], the decay in question receives a little less soft modification within the framework of LCSR, as expected. Also, it is easily figured out by intuitive picture that a bit more ratio is observed in the above two cases than $R = 0.30 - 0.70$ yielded for $B \to J/\psi K$ in [15].

To more give prominence to the role of the soft effects, we might, as in Eq(3), express $A = A_F + A_S$ as

$$A(\bar{B}^0 \to D^{*0}\pi^0) = -\frac{1}{2}G_FV_{cb}V_{ud}^*(m_b^2 - m_{D^*}^2)f_{D^*}F_1^{B^0 \to \pi}(m_{D^*}^2)a_2^{eff}, \quad (16)$$

in term of a phenomenological parameter $a_2^{eff} = a_2^f + 2C_2R_2$, with $a_2^f = C_1 + C_2/3$ corresponding to the factorizable contribution and $R_2$ being introduced to parameterize the soft effect. While $R_2$ turns out numerical smallness, the resulting influence on $a_2^{eff}$ is considerable, due to the large coefficient $2C_2$ multiplying it. Numerically, the soft effect is such that it adds a positive number to $a_2^f$, changing the value of $a_2^{eff}$ from 0.12 to 0.16 - 0.23. Such a result, along with that obtained in $\bar{B}^0 \to D^0\pi^0$, keeps still away from that demanded by the experiment data [3] $|a_2(\bar{B}^0 \to D^{0(*)}\pi^0)| = 0.57 \pm 0.06$.

At this point, a minor comment should be in order on a couple of open problems with soft dynamics in charmed color suppressed B decays: As a result of the dominance of nonfactorizable soft dynamics in $\bar{B}^0 \to D^{0(*)}\pi^0$, in addition to transverse momentum dependence which is important to get a result of infrared finiteness, the higher-twist components have to be counted in the light cone wavefunction description of $D^{0(*)}$ meson. The same must also be done for $B \to J/\psi K$, notwithstanding the fact that no corresponding infra divergence appears so as to validate QCD factorization for it, owing to the emitted quarkonium system $(C\bar{C})$ being of a small size $\sim 1/\alpha_s(m_c)m_c \ll 1/\Lambda_{QCD}$. On the other hand, as shown in $B \to \pi\pi, J/\psi K$, the hard spectator contributions have a divergent integral from the end point region of the distribution amplitude for the light mesons, when the twist-3 components is included. However, the transverse momenta may not be omitted in the region. We believe that for $\bar{B}^0 \to D^{0(*)}\pi^0$ it is this case too. The other resources of soft effects, of course, such as final state interactions [24], have to enter consideration in an effective manner for understanding the data.

Colligating the LCSR analyses made for $\bar{B}^0 \to D^{0(*)}\pi^0, J/\psi K$, we conclude that nonfactorizable higher twist corrections, in spite of being a nonleading effect, are important to understand the charmed color suppressed B decays. These findings favor,
in particular, the argument for QCD factorization: lange-distance QCD dynamics predominates in the color-suppressed processes such as $\bar{B}^0 \to D^{0(*)0}(\pi^0, \eta, \eta')$. A further investigation is necessary to improve the present results. The staple improvement is to take into account the QCD radiative correction to the correlator. The other mends are going to wait until the higher-twist amplitudes of pion become updated and available.
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Figure 1: Feynman diagram illustrating nonfactorizable higher twist contributions to the correlator (6). Solid lines represent quarks, dashed line gluon, wavy lines stand for external currents. The square refers to the insertion of the operator $\tilde{O}_1$, and oval the pionic higher twist wavefunctions. The cross denotes the other point to which gluon line can be attached.