**B → D K_{0,2}^* Decays:**

PQCD analysis to determine CP violation phase angle γ

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\[ B^\pm \to (D^0, D^0, D_{CP}) K_{0,2}^{\pm} \] decays are helpful in determining the CP violation angle γ, and we analyze these decay processes within the perturbative QCD approach based on \( k_T \) factorization. We found that the branching ratio of \( B^- \to D^0 K_{1}^- \) can reach the order of \( 10^{-4} \), due to the enhancement of nonfactorizable contributions in color-suppressed \( D^0 \)-emission, while the branching ratio of \( B^- \to \bar{D}^0 K_{0}^- \) is of the order \( 10^{-5} \). The ratio of decay amplitudes is about 3 times larger than the one in the channel \( B^\pm \to D K^\pm \). Large branching ratios provide a good opportunity to observe \( B^\pm \to D K_{0,2}^{\pm} \) on the ongoing and forthcoming experimental facilities and consequently these channels may be of valuable avail in reducing the errors in the CP violation phase angle γ. We also explore the possible time-dependent CP asymmetries of \( B_s \) decay into a scalar meson to determine the phase angle γ.

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**I. INTRODUCTION**

The authentication of the unitarity of CKM matrix allows us to explore the standard model (SM) description of the CP violation and reveal new physics beyond the SM. Among the angles (\( \alpha, \beta, \gamma \)) of the so-called (bd) unitarity triangle derived from the \( V_{ud}V_{ub}^* + V_{cd}V_{cb}^* + V_{td}V_{tb}^* = 0 \), satisfying the constraint \( \alpha + \beta + \gamma = 180^\circ \), the angle γ is least constrained, with a precision of roughly \( 10^\circ \). This is one of the main sources of the current uncertainties in the apex of the unitary triangle [1,2].

One of the most efficient ways proposed in the literature to measure γ makes use of the two triangles formed by the six channels of \( B^\pm \to (D^0, D^0, D_{CP}) K_{0,2}^{\pm} \). The shape of the two triangles is controlled by two quantities

\[ r_B^K = |A(B^- \to \bar{D}^0 K^-_1)/A(B^- \to D^0 K^-_2)|, \]
\[ \delta_B^K = \arg [e^{i\gamma}A(B^- \to \bar{D}^0 K^-_1)/A(B^- \to D^0 K^-_2)], \]

where \( K_1 \) can be \( K_0 \) or \( K_{0,2} \). One of the most intriguing properties in this method is that it is independent of hadronic uncertainties, and moreover the CP violation from the D meson decays can also be incorporated [3]. Due to the fact that the \( B^- \to \bar{D}^0 K^-_1 \) is both Cabibbo-suppressed and color suppressed, the ratio \( r_B^K \sim |V_{ub}V_{cs}^*/(V_{cb}V_{us})|a_2/a_1 | \sim 0.1 \) and in particular the world averages for these parameters [23]

\[ r_B^K = 0.107 \pm 0.010, \quad \delta_B^K = (112^{+12}_{-13})^\circ \]

indicate that the two triangles formed by decay amplitudes are squashed. As a consequence, the measurement of γ requests a precise knowledge on the \( B^- \to \bar{D}^0 K^-_1 \).

In Ref. [8], we proposed a new method to determine the CP violation angle γ that uses the \( B^\pm \to D K_{0,2}^* \) decays (see also Ref. [4]). Unlike the \( B \to D K^\pm \), the color-allowed amplitudes in \( B^\pm \to D K_{0,2}^* \) have vanishing/small decay constants and are comparable with the color-suppressed ones. Large interference between the two amplitudes is induced in the \( B \to D_{CP} K \) and the sensitivity to γ is greatly improved. Branching ratios of these channels are estimated to lie in the range from \( 10^{-6} \) to \( 10^{-5} \), using a method of factorization in conjunction with experimental data [8]. The motif of this work is to adopt the QCD-based factorization method, more explicitly the perturbative QCD (PQCD) approach [10,11] (see Ref. [12] and Ref. [14] for the recent developments and applications of the PQCD approach), to calculate the branching ratios, strong phases and CP asymmetries. The perturbative QCD approach is formulated on the basis of \( k_T \) factorization, and has been applied to B meson decays into charmed meson in a number of references and a global agreement of the results with the available data is found [13,22]. One of the most successful predictions is \( r_k = 0.092^{+0.012+0.003}_{-0.003-0.003} \), in good agreement with the data [22]. To the end of this work, we show that the resulting branching ratios are enhanced by one order of magnitude than our previous estimates, due to the inclusion of the large nonfactorizable contribution in \( D^0 \)-emission diagram. Such large branching ratios provide a better opportunity for the measurement of these channels on the experimental facilities and constraining the γ angle.

The rest of this work is organized as follows: In Sec. II, we will calculate the \( B \to D^0(\bar{D}^0)K_{0,2}^*(1430) \) decay amplitudes and give the factorization formulas, while Sec. III contains the numerical analysis and discussions. The last section is our summary. We also relegate some of the calculation details to the Appendix.
II. PERTURBATIVE QCD CALCULATION

In the PQCD approach, the inclusion of the intrinsic transverse momentum of valence quarks smears the endpoint singularities appearing in the calculations under the collinear factorization context. For the \( mb \to \infty \) limit, the decay amplitude is generically expressed as a convolution of wave functions and hard scattering kernel with both longitudinal momenta and transverse space coordinates

\[
\mathcal{M} = \int_0^1 dx_1 dx_2 dx_3 \int d^2 \vec{b}_1 d^2 \vec{b}_2 d^2 \vec{b}_3 \phi_B(x_1, \vec{b}_1, t) \\
\times T_H(x_1, x_2, \vec{b}_1, \vec{b}_2, t) \phi_2(x_2, \vec{b}_2, t) \phi_3(x_3, \vec{b}_3, t),
\]

where the \( B \) in the indices represents a \( B \) meson and 2, 3 represent the two mesons in the final state. In the computation of higher order QCD corrections, the overlap of soft and collinear momentum results in double logarithm divergences. Resummation of these leads to the Sudakov factor which has the tendency to diminish the endpoint contributions and supports the hard-scattering picture used in this framework. For a review of this approach, see Ref. [23].

The wave functions, the most important entry in the perturbative QCD approach, are nonperturbative in nature and can only be acquired by some nonperturbative methods or with the aid from some simple but effective models. For the \( B \) meson which is a heavy-light system, we adopt the light cone matrix

\[
\Phi_B = \frac{i}{\sqrt{2N_c}} (\bar{q} B + m_B) \gamma_5 \phi_B(x_1, \vec{b}_1),
\]

in which we have neglected the numerically-suppressed distribution amplitude [24]. Here \( x_1 \) is the momentum fraction of the light spectator quark and \( N_c = 3 \) is the color factor. As for the wave functions for the \( D \) meson, we use the form derived in Ref. [15]

\[
\Phi_D = \frac{i}{\sqrt{2N_c}} \gamma_5 (\bar{q} D + m_D) \phi_D(x_2, \vec{b}_2)
\]

The light-cone distribution amplitudes (LCDAs) for \( K_0^* \) are governed by the conformal spin symmetry of QCD and have the following definitions [25]

\[
\langle K_0^*(p_{K_0^*}) | q(0) \rangle \langle \bar{q}(z) | 0 \rangle = \frac{-1}{\sqrt{2N_c}} \int_0^1 dx_1 e^{ixp_{K_0^*} \cdot z} \\
\{ \bar{q} \phi_{K_0^*}(x) + m_{K_0^*} \phi'_{K_0^*}(x) + m_{K_0^*} (\bar{q} \vec{\gamma} - 1) \phi_{T_{K_0^*}}(x) \} j_{i},
\]

in which \( \bar{n} \) is chosen as the flight direction of the \( K_0^* \) in the \( B \) meson rest frame and \( n \) is opposite to \( \bar{n} \). These LCDAs, the twist-2 \( \phi_{K_0^*} \) and the twist-3 \( \phi'_{K_0^*}, \phi_{T_{K_0^*}} \), can be expanded in terms of Gegenbauer polynomials

\[
\phi_{K_0^*}(x) = \frac{f_{K_0^*}}{2\sqrt{2N_c}} 6x(1-x) \sum_{m=0}^\infty B_m c_m^{3/2}(2x-1),
\]

\[
\phi'_{K_0^*}(x) = \frac{f_{K_0^*}}{2\sqrt{2N_c}} \phi_{T_{K_0^*}}(x) = \frac{f_{K_0^*}}{2\sqrt{2N_c}} (1-2x),
\]

with \( B_m = (m_s - m_u)/m_{K^*} \). The decay constant \( f_{K_0^*} \) is defined by a scalar current

\[
\langle K_0^*(1430) | \bar{s} u | 0 \rangle = f_{K_0^*},
\]

and is related to the vector decay constant by \( f_{K_0^*} = B_0 f_{K_0^*} \). We will leave out the higher Gegenbauer moments in twist-3 LCDAs [26, 27] since their contributions are found to be typically small [28].

Similarly the LCDAs of a longitudinally polarized \( K_0^* \) state are defined as [29]

\[
\langle K_0^*(p_{K_0^*}, \epsilon) | q_{2\beta}(z) q_{1\alpha}(0) | 0 \rangle = \frac{1}{\sqrt{2N_c}} \int_0^1 dx_1 e^{ixp_{K_0^*} \cdot z} \\
\times \left\{ m_{K_0^*} \epsilon_{\alpha \beta} \phi_{K_0^*}(x) + m_{K_0^*} \epsilon_{\alpha \beta} \bar{p}_{K_0^*} \phi'_{K_0^*}(x) \right\} \delta_{\alpha \beta},
\]

with \( n^2 = v^2 = 0 \) being light-like unit vectors. The new vector \( \epsilon_{\alpha \beta} \) in Eq. 3 is related to the polarization tensor by \( \epsilon_{\alpha \beta} = \epsilon_{\alpha \beta} v \gamma \) and can be simplified in terms of a polarization vector

\[
\epsilon_{\alpha \beta} = \frac{\epsilon_{\alpha \beta} v \gamma}{p_{K_0^*} \cdot v} m_{K_0^*} \equiv \frac{\sqrt{2}}{\sqrt{3}} \epsilon_{\alpha \beta}.
\]

The above LCDAs have the asymptotic forms [29]

\[
\phi_{K_0^*}(x) = \frac{f_{K_0^*}}{2\sqrt{2N_c}} 30x(1-x)(2x-1),
\]

\[
\phi'_{K_0^*}(x) = \frac{f_{K_0^*}}{2\sqrt{2N_c}} \frac{15}{2} (2x-1)(1-6x+6x^2),
\]

\[
\phi_{T_{K_0^*}}(x) = \frac{f_{K_0^*}}{4\sqrt{2N_c}} d \int [15x(1-x)(2x-1)].
\]

There are three types of diagrams contributing to the decay amplitudes which are depicted in Fig. 1 [1] the color-allowed contributions in the process \( B^- \to D^0 K_{0*}^{*+}(1430)(a, b, c, d) \), the color-suppressed one in the process \( B^- \to D^0 K_{0*}^{*+}(1430)(e, f, g, h) \), and the annihilation one in the process \( B^- \to D^0 K_{0*}^{*+}(1430)(i, j, k, l) \). In the middle diagrams, the exchange of \( c \) and \( u \) quark results in the corresponding diagrams for the process \( B^- \to D^0 K_{0*}^{*+}(1430) \).

Factorization formulas for the \( K_0^* \)-emission diagrams are given as

\[
\xi_{ex} = N_1 f_{K_0^*} \int_0^1 dx_1 dx_3 \int_0^\infty b_1 db_1 b_3 d\beta \phi_B(x_1, \vec{b}_1) \\
\times \phi_D(x_3, \vec{b}_3) \left\{ (2 - x_3 + r_D(2x_3 - 1)) E_a(a_1) a_1(t_a) h_a + r_D(1 + r_D) E_b(t_b) a_1(t_b) h_b \right\},
\]

with \( N_1 = 8\pi C_F m_b^2 \). The hard kernels \( E_i(t_i) \) and \( h_i \) in these formulas are determined by the virtualities of the intermediate quarks and gluons and they can be found in
The nonfactorizable contributions, the last diagrams in Fig. 1, have the formulas

\[ M_{in} = N_2 \int_0^1 \int_0^{\infty} b_1 b_2 b_2 \phi_B(x_1, b_1) \phi_D(\bar{x}_2, b_2) \times \left\{ \begin{array}{l}
[2x_2 \phi_{K_1^0}(x_3) + r_{K^0}^2 \bar{x}_3 \phi_{K_1^0}(x_3) + \phi_{T_k^0}(x_3)] \\
h_T E_g(t_1) C_1(t_1) - h_T E_{h}(t_1) C_2(t_1) \\
\left[ (\bar{x}_2 + \bar{x}_3) \phi_{K_1^0}(x_3) + r_3 \bar{x}_3 \phi_{K_1^0}(x_3) - \phi_{T_k^0}(x_3) \right] \end{array} \right\} \]

while the \( D^0 \)-emission is factorized as

\[ M_{in}' = N_2 \int_0^1 \int_0^{\infty} b_1 b_2 b_2 \phi_B(x_1, b_1) \phi_D(x_2, b_2) \times \left\{ \begin{array}{l}
[2x_2 \phi_{K_1^0}(x_3) + r_{K^0}^2 \bar{x}_3 \phi_{K_1^0}(x_3) + \phi_{T_k^0}(x_3)] \\
h_T E'_g(t_1) C_2(t_1) - h_T E_{h}(t_1) C_2(t_1) \\
\left[ (\bar{x}_2 + \bar{x}_3) \phi_{K_1^0}(x_3) + r_3 \bar{x}_3 \phi_{K_1^0}(x_3) - \phi_{T_k^0}(x_3) \right] \end{array} \right\} \]

For the \( B \to D K_0^* \) decays, there are contributions from the annihilation diagrams, which are depicted in Fig. 1 (i)(j)(k)(l). The amplitude of factorizable annihilation diagrams are given as

\[ \xi_{exc} = N_1 f_B \int_0^1 dx_2dx_3 \int_0^{\infty} b_2 b_3 b_3 \phi_D(x_3, b_3) \times \left\{ \begin{array}{l}
[2(x_2 + 1)] r_D r_{K_0^*} \phi_{K_0^*}(x_2) \\
-x_3 \phi_{K_0^*}(x_2)] + E_l(t_1) h_1 a_1(t_1) \\
\left[ r_D r_{K_0^*}(2x_2 - 3) \phi_{K_0^*}^s(x_2) - (2x_2 - 1) \phi_{K_0^*}^T(x_2) \right] \\
-(x_2 - 1) \phi_{K_0^*}(x_2) \end{array} \right\} \]

The nonfactorizable annihilation amplitude is given by

\[ M_{exc} = N_2 \int_0^1 dx_2dx_3 \int_0^{\infty} b_2 b_3 b_3 \phi_D(x_3, b_3) \times \left\{ \begin{array}{l}
[2(x_2 + 1)] r_D r_{K_0^*}(x_2) \\
-x_3 \phi_{K_0^*}(x_2)] + E_l(t_1) h_1 a_1(t_1) \\
\left[ r_D r_{K_0^*}(2x_2 - 3) \phi_{K_0^*}^s(x_2) - (2x_2 - 1) \phi_{K_0^*}^T(x_2) \right] \\
+x_3 \phi_{K_0^*}(x_2) \end{array} \right\} \]

The formulas for channels involving \( K_2^* \) are obtained by the replacement \( f_{K_0^*} \to 0, \phi_{K_0^*} \to \phi_{K_2^*} \) and \( \phi_{K_2^*} \to \phi_{K_2^*}^{T*} \). Incorporating the CKM matrix elements, we have the total decay amplitudes

\[ A(B^- \to D^0 K_0^*) = \frac{G_F}{\sqrt{2}} V_{ub} V_{cs}^* \times \left( \xi_{in} + M_{in} + \xi_{exc} + M_{exc} \right), \]

\[ A(B^- \to D^0 K_0^{*+}) = \frac{G_F}{\sqrt{2}} V_{ub} V_{us}^* \times \left( \xi_{exc} + M_{exc} + \xi_{in} + M_{in} \right). \]
where $G_F$ is the Fermi constant. It should be pointed out that the color-allowed $\xi_q$ is zero in $B^− \to D^0 K^−$ due to the fact that the tensor meson can not be generated by a local vector or axial-vector current.

III. NUMERICAL RESULTS AND DISCUSSIONS

The expression for $\phi_B(x, b)$ has been examined in various kinds of $B$ decays and the currently-accepted form in the PQCD approach is

$$\phi_B(x, b) = N_B x^2 (1-x)^2 \exp \left[ -m_{b^+}^2 x^2 - \frac{1}{2} (\omega x b)^2 \right], \quad (13)$$

where the normalization factor $N_B$ is related to the decay constant $f_B$. We adopt the ansatz that $B$ meson wave functions have a sharp cone at $x \approx 0.1$, in accordance with the most probable momentum fraction of the light quark. The best-fitted form for $\phi_D$ from the $B$ meson decays into a charmed meson derived in Refs. \textit{[16–21]} is

$$\phi_D(x_2, b_2) = \frac{f_D}{2\sqrt{2}N} 6x(1-x)\left[1+C_D(1-2x)\right] \times \exp \left[-\omega_D x^2 b^2/2\right]. \quad (14)$$

Their numerical values (in GeV except $C_D$) are used as

$$C_D = (0.5 \pm 0.1), \quad \omega_3 = (0.40 \pm 0.05), \quad \omega_D = 0.1,$$

$$f_B = (0.1969 \pm 0.0089), \quad f_D = (0.221 \pm 0.018), \quad (15)$$

where $f_B$ is from the recent Lattice QCD simulation \textit{[30]} and the $f_D$ is extracted from $D^- \to \mu \bar{\nu}_{\mu}$ \textit{[31]}. For the LCDAs of the light scalar meson $K^*_0$, we adopt $B_0 = (m_s - m_u)/m_K = 0.07$ \textit{[31]} and the two different solutions in Ref. \textit{[22–23]}

$$S1 : \quad \tilde{f}_{K^*_0} = (-300 \pm 30)\text{MeV}, \quad B_1 = 0.58 \pm 0.07, \quad B_3 = -1.20 \pm 0.08,$$

$$S2 : \quad \tilde{f}_{K^*_0} = (445 \pm 50)\text{MeV}, \quad B_1 = -0.57 \pm 0.13, \quad B_3 = -0.42 \pm 0.22. \quad (16)$$

The normalization constants in $K^*_2$ LCDAs are \textit{[24]}

$$f_{K^*_2} = (118 \pm 5)\text{MeV}, \quad f_{K^*_2} = (77 \pm 14)\text{MeV}. \quad (17)$$

These LCDAs have been used to calculate the form factors of $B$ decays into a scalar tensor meson in the same perturbative QCD approach \textit{[32–35]}. For the CKM matrix elements, we use \textit{[31]}

$$|V_{ub}| = (3.89 \pm 0.44) \times 10^{-3},$$

$$|V_{cs}| = 0.97345,$$

$$|V_{us}| = 0.2252,$$

$$|V_{ub}| = (40.6 \pm 1.3) \times 10^{-3}, \quad (18)$$

where the small uncertainties are not taken into account.

With the above inputs, we predict the branching ratios as

$$B(B^− \to D^0 K^+_1) = (2.70^{+1.09+0.25}_{−0.97−0.48}) \times 10^{-4}, \quad S1$$

$$B(B^− \to D^0 K^+_2) = (1.53^{+0.82+0.62}_{−0.53−0.46}) \times 10^{-5}, \quad S1$$

$$B(B^− \to D^0 K^+_3) = (1.61^{+0.49+0.44}_{−0.41−0.22}) \times 10^{-4}, \quad S2$$

$$B(B^− \to D^0 K^+_4) = (3.38^{+1.51+0.65}_{−1.18−0.59}) \times 10^{-5}, \quad S2$$

$$B(B^− \to D^0 K^+_5) = (2.40^{+1.30+0.72}_{−0.97−1.05}) \times 10^{-5}, \quad S2$$

$$B(B^− \to D^0 K^+_6) = (3.32^{+1.90+0.77}_{−1.88−0.74}) \times 10^{-6}, \quad (19)$$

where the first uncertainties are from $f_B$ and $\omega_b$ in the $B$ meson wave functions, the second errors are from $A_{QCD}$ and the scales defined in Appendix A (We vary the $\sqrt{Q}$ and $\sqrt{P}$ in the scales 25% for error estimation). The results for the branching ratios made here are larger than our previous estimates in Ref. \textit{[3]}, obtained under the factorization approach. The main reason is due to the enhancement of nonfactorizable contributions in color-suppressed $D^0$-emission. Ratios and phases of the amplitudes are

$$r_{K^*_0} = 0.24^{+0.02+0.07}_{−0.01−0.04} \quad \delta_{K^*_0} = (-125.65^{+3.95+23.16}_{−9.01−17.42})^\circ, \quad S1$$

$$r_{K^*_0} = 0.54^{+0.03+0.07}_{−0.02−0.04} \quad \delta_{K^*_0} = (-161.51^{+0.91+12.01}_{−0.96−9.53})^\circ, \quad S2$$

$$r_{K^*_2} = 0.37^{+0.02+0.17}_{−0.00−0.09} \quad \delta_{K^*_2} = (155.53^{+0.00+2.98}_{−3.36−3.49})^\circ. \quad (20)$$

It should be pointed out that although the large uncertainties in many entries like decay constants will affect our predictions for branching ratios, the relative strength of decay amplitudes are almost unaffected.

Physical observables that are experimentally explored are defined as

$$R_{CP}^{K^+} = 2 \frac{B(B^− \to D_{CP}^0 K^−) + B(B^+ \to D_{CP}^+ K^+_1)}{B(B^− \to D^0 K^+_1) + B(B^+ \to D^0 K^+_1)}$$

$$= 1 + (r_{K^*_0}^2 B)^2 + 2 r_{K^*_0} \cos \delta_{K^*_0} \gamma_{RCP},$$

$$A_{CP}^{K^+} = \frac{B(B^− \to D_{CP}^0 K^−) - B(B^+ \to D_{CP}^+ K^+_1)}{B(B^− \to D^0 K^+_1) + B(B^+ \to D^0 K^+_1)}$$

$$= \pm 2 r_{K^*_0} \sin \delta_{K^*_0} \gamma_{RCP}. \quad (21)$$

In the limit of $r_B \to 0$, the ratio $R_{CP}^{K^+}$ is close to 1 while the CP asymmetries vanish. As we have pointed out, due to the suppression of the color-allowed decay amplitudes based on the fact that the matrix element of a local vector or axial-vector current (at the lowest order in $\alpha_s$) between the QCD vacuum and the $K^*_0$ ($K^*_2$) state is small (identically zero), the low sensitivity to $\gamma$ is improved and in particular large CP asymmetries are expected. The dependence of $R_{CP}$ and $A_{CP}$ on $\gamma$ is shown in Fig. 2. Since the errors of $r_{K^*_0}$ and $\delta_{K^*_0}$ are not large, only their central values are used. We investigate these observables in the region $\gamma = (68^{+10}_{−5})^\circ$ which is from a combined analysis of $B^\pm \to D K^\pm$. In this region we find that the observables of the $B^− \to (D^0, D^0) K^0_2$ in $S1$ have relative smaller variances because of the smaller $r_{K^*_0}$, most of which are around 10%. However, for the other cases the
observables have large variances, and some of them even reach about 40%. Therefore these channels have the potential to improve the accuracy of $\gamma$ extracted from the $B^\pm \to DK^\pm$ decays.

It is also interesting to notice that due to the large value of $r_{K_{0,2}^*}$, the large impact arising from the direct CP violation of $D^0$ decays into $CP$ eigenstates $K^+K^-/\pi^+\pi^-$, of the order $\mathcal{O}(A_{CP}^2)/r_{K_{0,2}^*}$ are not important in $B \to DK_{0,2}^*$.

As discussed in Ref. [8], the time-dependent observables of $B_s \to (D, \bar{D})f_0(980)$ and $B_s \to (D, \bar{D})f_2'(1525)$ processes can be used to determine the $\gamma$ as well. Therefore we will also predict their branching ratios in the perturbative QCD approach. In these channels only the color-suppressed diagrams depicted as Fig. 1(c) contribute, in which the spectator quark $\bar{u}$ need be replaced by $\bar{s}$ and the $c$ and $\bar{u}$ in the emission meson should be exchanged for the $B_s \to D(f_0, f_2')$ decays. The amplitudes are given by

$$A(\bar{B}_s^0 \to D^0 f_0, f_2') = \frac{G_F}{\sqrt{2}} V_{ub} V_{cs}^* (\xi'_u + \mathcal{M}_{in}),$$
$$A(\bar{B}_s^0 \to D^0 f_0, f_2') = \frac{G_F}{\sqrt{2}} V_{ub} V_{cs}^* (\xi'_c + \mathcal{M}_{in}). \quad (22)$$

Our inputs for the $\bar{B}_s^0 \to (\bar{D}, D)(f_0, f_2')$ decays are summarized as (decay constants in units of GeV) [25, 29, 30]

$$f_B = 0.2420 \pm 0.0005, \quad f_{f_0} = 0.37 \pm 0.02,$$
$$B_1(f_0) = -0.78 \pm 0.08, \quad B_1(f_0) = 0.02 \pm 0.07,$$
$$f_{f_2'} = 0.126 \pm 0.004, \quad f_{f_2'} = 0.065 \pm 0.012, \quad (23)$$

with which the branching ratios are predicted as

$$B(\bar{B}_s^0 \to D^0 f_0) = (3.50^{+1.26+0.56}_{-1.15-0.77}) \times 10^{-5},$$
$$B(\bar{B}_s^0 \to D^0 f_0) = (5.94^{+3.13+1.47}_{-3.16-0.97}) \times 10^{-6},$$
$$B(\bar{B}_s^0 \to D^0 f_2') = (1.08^{+0.37+0.26}_{-0.35-0.28}) \times 10^{-5},$$
$$B(\bar{B}_s^0 \to D^0 f_2') = (2.85^{+1.53+0.67}_{-0.95-0.43}) \times 10^{-6}. \quad (24)$$

IV. SUMMARY

The determination of the CKM angles is crucial for the test of the CKM paradigm and also sheds light on the standard model description of the CP violation. To accomplish this goal, one of the most important efforts to be done in the next step is to reduce the uncertainties in these entries. What has been explored in Ref. [8] and this work is to propose that the $B \to DK_{0,2}^*$ is expedient to provide complementary information of the angle $\gamma$.

In this work we have calculated the branching ratios of $B \to DK_{0,2}^*$ and the corresponding $B_s$ relatives, by adopting the $k_T$ factorization approach. We find that the BR of $B \to DK_{0,2}^*$ can reach $10^{-4}$ ($10^{-5}$) while the ratio of the magnitude is also significantly enhanced compared to $B \to DK$ mode. As a consequence, it seems promising for the LHCb experiment and the currently-designed Super B factory to measure $B_{u,d} \to DK_{0,2}(1430)$ and time-dependent CP asymmetries in the $B_s$ decays.

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Appendix A: Hard kernels in the PQCD calculation

The offshellness of the intermediate gluon

\[ Q_{a,b,c,d} = x_1 \bar{x}_3 m_B^2, \quad Q_{e,f,g,h',h''} = x_1 \bar{x}_3 (1 - r_D^2) m_B^2, \]

\[ Q_{k,l,m,n} = x_3 \bar{x}_2 (1 - r_D^2) m_B^2, \]

and the quarks

\[ P_a = \bar{x}_3 m_B^2, \quad P_b = x_1 m_B^2, \]

\[ P_c = \bar{x}_3 (x_1 - x_2) (1 - r_D^2) m_B^2, \]

\[ P_d = \bar{x}_3 (x_1 - x_2) (1 - r_D^2) m_B^2, \]

\[ P_e = \{ (x_1 - x_2) [ (1 - r_D^2) \bar{x}_3 + r_D^2 ] + r_D^2 \} m_B^2, \]

\[ P_g = \{ (x_1 - x_2) [(1 - r_D^2) \bar{x}_3 + r_D^2 ] + r_D^2 \} m_B^2, \]

\[ P_h = \{ (x_1 - x_2) [(1 - r_D^2) \bar{x}_3 + r_D^2 ] + r_D^2 \} m_B^2, \]

\[ P_k = (x_1 - x_2) (1 - r_D^2) m_B^2, \]

\[ P_m = -(\bar{x}_3 [1 - (1 - r_D^2) \bar{x}_2 - x_1] - 1] m_B^2, \]

\[ P_n = x_3 [x_1 - \bar{x}_2 (1 - r_D^2)] m_B^2, \]

result in the hard scales and kernels

\[ t_i = \max \{ \sqrt{Q_1}, \sqrt{P_1}, 1/b_1, 1/b_3 \}, \]

\[ t_j = \max \{ \sqrt{Q_2}, \sqrt{P_2}, 1/b_1, 1/b_3 \}, \]

\[ t_x = \max \{ \sqrt{Q_3}, \sqrt{P_3}, 1/b_1, 1/b_3 \}, \]

\[ t_y = \max \{ \sqrt{Q_4}, \sqrt{P_4}, 1/b_1, 1/b_3 \}, \]

\[ h_{a,c} = H_{a,c} (P_{a,c}, Q_{a,c}, b_1, b_3) S_i (x_3), \]

\[ h_{b,f} = H_{b,f} (P_{b,f}, Q_{b,f}, b_3, b_1) S_i (x_3), \]

\[ h_j = H_{j} (Q_j, P_j, b_1, b_2), \]

\[ h_k = H_{a,f} (P_k, Q_k, b_2, b_3) S_i (x_3), \]

\[ h_l = H_{l} (P_l, Q_l, b_3, b_2) S_i (x_2), \]

\[ h_{m,n} = H_{m,n}, \quad (A1) \]

for the factorizable diagram \( i = a, b, c, f, \) for the nonfactorizable diagram \( j = c, d, g, h', h'', \) and for the annihilation diagrams \( x = k, l \) and \( y = m, n. \) For diagrams \( (c,d), \) we also keep the function \( S_i (x_3). \) Here we use the definition of Bessel functions

\[ H_{c} (\alpha, \beta, b_1, b_3) = K_0 (\sqrt{\beta b_1}) [\theta (b_1 - b_3) J_0 (\sqrt{\alpha b_3}) \]

\[ \times K_0 (\sqrt{\alpha b_1}) + (b_1 \leftrightarrow b_3)], \]

\[ H_{cn} (\alpha, \beta, b_1, b_2) = [\theta (b_2 - b_1) K_0 (\sqrt{\alpha b_2}) \times I_0 (\sqrt{\alpha b_1}) \]

\[ + (b_1 \leftrightarrow b_2)] \left\{ \begin{array}{ll} \frac{\pi}{2} H_0^{(1)} (\sqrt{\beta b_1}), & \beta < 0 \\ K_0 (\sqrt{\beta b_1}), & \beta > 0 \end{array} \right. \quad (A2) \]

with the quark anomalous dimension \( \gamma_q = -\alpha_s / \pi. \)
