PHYSICAL REVIEW D **88,** 034003 (2013)

$B \to DK^*_{0,2}$ decays and the CP phase angle γ

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(Received 28 May 2013; published 2 August 2013)

$B^\pm \to (D^0, \bar{D}^0, D_{CP})K^{*\pm}_{0,2}$ 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 $B^- \to D^0 K_0^{*-}$ 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 $B^\pm \to DK^\pm$. Large branching ratios provide a good opportunity to observe $B^\pm \to DK_2^{*\pm}$ 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.

DOI: 10.1103/PhysRevD.88.034003 PACS numbers: 13.25.Hw, 12.15.Hh

I. INTRODUCTION

The authentication of the unitarity of the Cabibbo-Kobayashi-Maskawa (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 (α, β, γ) 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 $\alpha + \beta + \gamma = 180°$, the angle γ is least constrained, with a precision of roughly 10°. 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 $B^\pm \to (D^0, \bar{D}^0, D_{CP})K^\pm$ [3–5]. The shape of the two triangles is controlled by two quantities,

$$r_B^{K_J} \equiv |A(B^- \to \bar{D}^0 K_J^-)/A(B^- \to D^0 K_J^-)|,$$
$$\delta_B^{K_J} \equiv \arg[e^{i\gamma}A(B^- \to \bar{D}^0 K_J^-)/A(B^- \to D^0 K_J^-)],$$

where K_J can be K 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 [6]. Due to the fact that the $B^- \to \bar{D}^0 K^-$ is both Cabibbo suppressed and color suppressed, the ratio $r_B^K \sim |V_{ub}V_{cs}^*/(V_{cb}V_{us}^*)a_2/a_1| \sim 0.1$ is small and in particular the world averages for these parameters [7],

$$r_B^K = 0.107 \pm 0.010, \qquad \delta_B^K = (112^{+12}_{-13})°,$$

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

In Ref. [8], we proposed a new method to determine the CP violation angle γ that uses the $B^\pm \to DK^{*\pm}_{0,2}$ decays (see also Ref. [9]). Unlike the $B \to DK^\pm$, the color-allowed amplitudes in $B^\pm \to DK^{*\pm}_{0,2}$ 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 [10–12] (see Refs. [13,14] 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 [15–22]. One of the most successful predictions is $r_K = 0.092^{+0.012+0.003}_{-0.003-0.003}$ [21], in good agreement with the data [7]. To the end of this work, we show that the resulting branching ratios are enhanced by 1 order of magnitude higher 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.

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II. PERTURBATIVE QCD CALCULATION

In the PQCD approach, the inclusion of the intrinsic transverse momentum of valence quarks smears the end point singularities appearing in the calculations under the collinear factorization context. In the $m_b \to \infty$ 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,

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

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 them leads to the Sudakov factor which has the tendency to diminish the end point 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,

$$\Phi_B = \frac{i}{\sqrt{2N_c}}(\not{p}_B + m_B)\gamma_5\phi_B(x_1, b_1), \quad (2)$$

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]:

$$\Phi_D = \frac{i}{\sqrt{2N_c}}\gamma_5(\not{p}_D + m_D)\phi_D(x_2, b_2). \quad (3)$$

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

$$\langle K_0^*(p_{K_0^*})|q(0)_j\bar{q}(z)_l|0\rangle = \frac{-1}{\sqrt{2N_c}}\int_0^1 dx e^{ixp_{K_0^*}\cdot z}\{\not{p}_{K_0^*}\phi_{K_0^*}(x) + m_{K_0^*}\phi^s_{K_0^*}(x) + m_{K_0^*}(\not{n}\not{\bar{n}} - 1)\phi^T_{K_0^*}(x)\}_{jl},$$

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 $\phi^{s,T}_{K_0^*}$, can be expanded in terms of Gegenbauer polynomials,

$$\phi_{K_0^*}(x) = \frac{\bar{f}_{K_0^*}}{2\sqrt{2N_c}}6x(1-x)\sum_{m=0}^{\infty} B_m C_m^{3/2}(2x-1),$$
$$\phi^s_{K_0^*}(x) = \frac{\bar{f}_{K_0^*}}{2\sqrt{2N_c}}, \qquad \phi^T_{K_0^*}(x) = \frac{\bar{f}_{K_0^*}}{2\sqrt{2N_c}}(1-2x), \quad (4)$$

with $B_0 = (m_s - m_u)/m_{K_0^*}$. The decay constant $\bar{f}_{K_0^*}$ is defined by a scalar current,

$$\langle K_0^{*-}(1430)|\bar{s}u|0\rangle = \bar{f}_{K_0^*},$$

and is related to the vector decay constant by $f_{K_0^*} = B_0\bar{f}_{K_0^*}$. 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_2^* state are defined as [29]

$$\langle K_2^*(p_{K_2^*}, \epsilon)|\bar{q}_{2\beta}(z)q_{1\alpha}(0)|0\rangle = \frac{1}{\sqrt{2N_c}}\int_0^1 dx e^{ixp_{K_2^*}\cdot z}\{m_{K_2^*}\not{\epsilon}^*_{\bullet L}\phi_{K_2^*}(x) + \not{\epsilon}^*_{\bullet L}\not{p}_{K_2^*}\phi^t_{K_2^*}(x) + m^2_{K_2^*}\frac{\epsilon_{\bullet L}\cdot v}{p_{K_0^*}\cdot v}\phi^s_{K_2^*}(x)\}_{\alpha\beta}, \quad (5)$$

with $n^2 = v^2 = 0$ being lightlike unit vectors. The new vector $\epsilon_{\bullet L}$ in Eq. (5) is related to the polarization tensor by $\epsilon_{\bullet L\mu} \equiv \frac{\epsilon_{\mu\nu}v^\nu}{p_{K_2^*}\cdot v}m_{K_2^*}$ and can be simplified in terms of a polarization vector,

$$\epsilon_{\bullet L\mu} \equiv \frac{\epsilon_{\mu\nu}v^\nu}{p_{K_2^*}\cdot v}m_{K_2^*} \simeq \frac{\sqrt{2}}{\sqrt{3}}\epsilon_{L\mu}. \quad (6)$$

The above LCDAs have the asymptotic forms [29],

$$\phi_{K_2^*}(x) = \frac{f_{K_2^*}}{2\sqrt{2N_c}}30x(1-x)(2x-1),$$
$$\phi^t_{K_2^*}(x) = \frac{f^T_{K_2^*}}{2\sqrt{2N_c}}\frac{15}{2}(2x-1)(1-6x+6x^2), \quad (7)$$
$$\phi^s_{K_2^*}(x) = \frac{f^T_{K_2^*}}{4\sqrt{2N_c}}\frac{d}{dx}[15x(1-x)(2x-1)].$$

There are three types of diagrams contributing to the decay amplitudes which are depicted in Fig. 1: the color-allowed contributions in the process $B^- \to D^0K^{*-}_{0(2)}(1430)(a, b, c, d)$, the color-suppressed one in the process $B^- \to D^0K^{*-}_{0(2)}(1430)(e, f, g, h)$, and the annihilation one in the process $B^- \to \bar{D}^0K^{*-}_{0(2)}(1430)(i, j, k, l)$. In the middle four diagrams, the exchange of c and u quark results in the corresponding diagrams for the process $B^- \to \bar{D}^0K^{*-}_{0(2)}(1430)$.

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

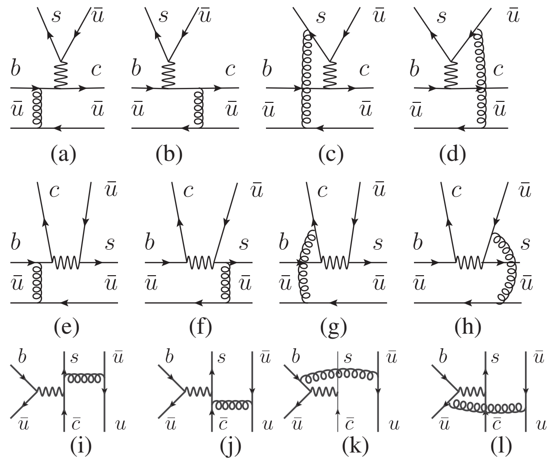

FIG. 1. Feynman diagrams for the color-allowed contributions in the process $B^- \to D^0 K^{*-}_{0(2)}(1430)$: (a), (b), (c), and (d); for the color-suppressed contributions in the process $B^- \to D^0 K^{*-}_{0(2)}(1430)$: (e), (f), (g), and (h); and for the annihilation contributions in process $B^- \to \bar{D}^0 K^{*-}_{0(2)}(1430)$: (i), (j), (k), and (l). In the middle four diagrams, the replacement of the $c\bar{u}$ by $u\bar{c}$ results in the corresponding diagrams for the process $B^- \to \bar{D}^0 K^{*-}_{0(2)}(1430)$.

$$\begin{aligned}\xi_{ex} = {} & N_1 f_{K_0^*} \int_0^1 dx_1 dx_3 \int_0^\infty b_1 db_1 b_3 db_3 \phi_B(x_1, b_1) \\ & \times \phi_D(\bar{x}_3, b_3)[(2 - x_3 + r_D(2x_3 - 1))E_a(t_a)a_1(t_a)h_a \\ & + r_D(1 + r_D)E_b(t_b)a_1(t_b)h_b],\end{aligned} \tag{8}$$

with $N_1 = 8\pi C_F m_B^4$. 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 Appendix. The nonfactorizable contributions, depicted by Figs. 1(c) and 1(d), have the formulas

$$\begin{aligned}\mathcal{M}_{ex} = {} & N_2 \int_0^1 [dx] \int_0^\infty b_1 db_1 b_2 db_2 \phi_B(x_1, b_1) \phi_D(\bar{x}_3, b_3) \\ & \times \phi_{K_0^*}(x_2)[(x_2 - r_D \bar{x}_3)E_c(t_c)h_c C_1(t_c) \\ & + (\bar{x}_3 r_D - \bar{x}_3 - \bar{x}_2)E_d(t_d)h_d C_1(t_d)],\end{aligned}$$

with $\bar{x}_i = 1 - x_i$, $N_2 = 32\pi m_B^4 C_F/\sqrt{2N_c}$, $[dx] \equiv dx_1 dx_2 dx_3$. In the collinear approximation, the amplitude is divergent when the momentum fraction of the light spectator in the final state goes to zero. The transverse momentum regulates this end point singularity and the threshold resummation function S_t will suppress the end point contribution further.

The factorizable color-suppressed diagrams, either the D^0 or $\bar{D}^0$ emission diagrams, have the same factorization formulas:

$$\begin{aligned}\xi_{\rm in} = {} & \xi'_{\rm in} \\ = {} & N_1 f_D \int_0^1 dx_1 dx_3 \int_0^\infty b_1 db_1 b_3 db_3 \phi_B(x_1, b_1) \\ & \times \{E_e(t_e)h_e a_2(t_e)[r_{K_0^*}(1 - 2x_3)(\phi^s_{K_0^*}(x_3) - \phi^T_{K_0^*}(x_3)) \\ & + (2 - x_3)\phi_{K_0^*}(x_3)] - 2r_{K_0^*}\phi^s_{K_0^*}(x_3)E_f(t_f)h_f a_2(t_f)\},\end{aligned} \tag{9}$$

but the D^0-emission nonfactorizable diagram is

$$\begin{aligned}\mathcal{M}_{\rm in} = {} & N_2 \int_0^1 [dx] \int_0^\infty b_1 db_1 b_2 db_2 \phi_B(x_1, b_1) \phi_D(\bar{x}_2, b_2) \\ & \times \{[x_2 \phi_{K_0^*}(x_3) + r_{K_0^*}\bar{x}_3(\phi^s_{K_0^*}(x_3) + \phi^T_{K_0^*}(x_3))] \\ & \times h_g E_g(t_g)C_1(t_h) - h_h E_h(t_h)C_2(t_h) \\ & \times [(\bar{x}_2 + \bar{x}_3)\phi_{K_0^*}(x_3) + r_3 \bar{x}_3(\phi^s_{K_0^*}(x_3) - \phi^T_{K_0^*}(x_3))]\},\end{aligned}$$

while the $\bar{D}^0$ emission is factorized as

$$\begin{aligned}\mathcal{M}'_{\rm in} = {} & N_2 \int_0^1 [dx] \int_0^\infty b_1 db_1 b_2 db_2 \phi_B(x_1, b_1) \phi_D(x_2, b_2) \\ & \times \{[x_2 \phi_{K_0^*}(x_3) + r_{K_0^*}\bar{x}_3(\phi^s_{K_0^*}(x_3) \\ & + \phi^T_{K_0^*}(x_3))]h'_g E'_g(t'_g)C_2(t'_g) \\ & - h'_h E'_h(t'_h)C_2(t'_h)[(\bar{x}_2 + \bar{x}_3)\phi_{K_0^*}(x_3) \\ & + r_3 \bar{x}_3(\phi^s_{K_0^*}(x_3) - \phi^T_{K_0^*}(x_3)]\}.\end{aligned} \tag{10}$$

For the $B \to \bar{D}K_0^*$ decays, there are contributions from the annihilation diagrams, which are depicted in Figs. 1(i)–1(l). The amplitude of factorizable annihilation diagrams is given as

$$\begin{aligned}\xi_{\rm exc} = {} & N_1 f_B \int_0^1 dx_2 dx_3 \int_0^\infty b_2 db_2 b_3 db_3 \phi_D(x_3, b_3) \\ & \times \{E_k(t_k)h_k a_1(t_k)[2(x_3 + 1)r_D r_{K_0^*}\phi^s_{K_0^*}(x_2) \\ & - x_3 \phi_{K_0^*}(x_2))] + E_l(t_b)h_l a_1(t_l) \\ & \times [r_D r_{K_0^*}((2x_2 - 3)\phi^s_{K_0^*}(x_2) \\ & - (2x_2 - 1)\phi^T_{K_0^*}(x_2)) - (x_2 - 1)\phi_{K_0^*}(x_2)]\}.\end{aligned} \tag{11}$$

The nonfactorizable annihilation amplitude is given by

$$\begin{aligned}\mathcal{M}_{\rm exc} = {} & N_2 \int_0^1 [dx] \int_0^\infty b_1 db_1 b_2 db_2 \phi_B(x_1, b_1) \phi_D(x_3, b_3) \\ & \times \{[r_D r_{K_0^*}((x_2 + x_3 - 1)\phi^T_{K_0^*}(x_2) \\ & - (x_2 - x_3 - 3)\phi^s_{K_0^*}(x_2)) \\ & + (x_2 - 1)\phi_{K_0^*}(x_2)]h_m E_m(t_m)C_1(t_m) \\ & + h_n E_n(t_n)C_1(t_n)[r_D r_{K_0^*}((x_2 - x_3 - 1)\phi^s_{K_0^*}(x_2) \\ & + (x_2 + x_3 - 1)\phi^T_{K_0^*}(x_2)) + x_3 \phi_{K_0^*}(x_2)]\}.\end{aligned}$$

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 $\phi_{K_0^*}^{s,T} \to \phi_{K_2^*}^{s,t}$. Incorporating the CKM matrix elements, we have the total decay amplitudes,

$$A(B^- \to \bar{D}^0 K_{0,2}^{*-}) = \frac{G_F}{\sqrt{2}} V_{ub} V_{cs}^* (\xi_{\rm in}' + \mathcal{M}_{\rm in}' + \xi_{\rm exc} + \mathcal{M}_{\rm exc}),$$
$$A(B^- \to D^0 K_{0,2}^{*-}) = \frac{G_F}{\sqrt{2}} V_{cb} V_{us}^* (\xi_{\rm ex} + \mathcal{M}_{\rm ex} + \xi_{\rm in} + \mathcal{M}_{\rm in}), \tag{12}$$

where G_F is the Fermi constant. It should be pointed out that the color-allowed $\xi_{\rm ex}$ is zero in $B^- \to D^0 K_2^{*-}$ due to the fact that the tensor meson cannot 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

$$\phi_B(x, b) = N_B x^2 (1-x)^2 \exp\left[-\frac{m_B^2 x^2}{2\omega_b^2} - \frac{1}{2}(\omega_b b)^2\right], \tag{13}$$

where the normalization factor N_B is related to the decay constant f_B. We adopt the ansatz that the B meson wave functions have a sharp peak at $x \sim 0.1$, in accordance with the most probable momentum fraction of the light quark. The best-fitted form for ϕ_D from the B meson decays into a charmed meson derived in Refs. [16,20,21] is

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

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

$$C_D = (0.5 \pm 0.1), \qquad \omega_b = (0.40 \pm 0.05),$$
$$\omega_D = 0.1, \qquad f_B = (0.1969 \pm 0.0089), \tag{15}$$
$$f_D = (0.221 \pm 0.018),$$

where f_B is from the recent lattice QCD simulation [30] and the f_D is extracted from $D^- \to \mu \bar{\nu}_\mu$ [31].

For the LCDAs of the light scalar meson K_0^*, we adopt $B_0 = (m_s - m_u)/m_{K_0^*} = 0.07$ [31] and the two different solutions in Ref. [25]:

$$S1: \bar{f}_{K_0^*} = (-300 \pm 30)\ \text{MeV}, \quad B_1 = 0.58 \pm 0.07, \quad B_3 = -1.20 \pm 0.08,$$
$$S2: \bar{f}_{K_0^*} = (445 \pm 50)\ \text{MeV}, \quad B_1 = -0.57 \pm 0.13, \quad B_3 = -0.42 \pm 0.22. \tag{16}$$

The normalization constants in K_2^* LCDAs are [29]

$$f_{K_2^*} = (118 \pm 5)\ \text{MeV}, \qquad f_{K_2^*}^T = (77 \pm 14)\ \text{MeV}. \tag{17}$$

These LCDAs have been used to calculate the form factors of B decays into a scalar/tensor meson in the same perturbative QCD approach [32–35].

For the CKM matrix elements, we use [31]

$$|V_{ub}| = (3.89 \pm 0.44) \times 10^{-3}, \qquad |V_{cs}| = 0.97345,$$
$$|V_{us}| = 0.2252, \qquad |V_{cb}| = (40.6 \pm 1.3) \times 10^{-3}, \tag{18}$$

where the small uncertainties are not taken into account.

With the above inputs, we predict the branching ratios as

$$\mathcal{B}(B^- \to D^0 K_0^{*-}) = (2.70^{+1.09+0.25}_{-0.97-0.48}) \times 10^{-4}, \quad S1$$
$$\mathcal{B}(B^- \to \bar{D}^0 K_0^{*-}) = (1.53^{+0.82+0.62}_{-0.53-0.46}) \times 10^{-5}, \quad S1$$
$$\mathcal{B}(B^- \to D^0 K_0^{*-}) = (1.16^{+0.40+0.14}_{-0.41-0.22}) \times 10^{-4}, \quad S2$$
$$\mathcal{B}(B^- \to \bar{D}^0 K_0^{*-}) = (3.38^{+1.51+0.65}_{-1.18-0.59}) \times 10^{-5}, \quad S2 \tag{19}$$
$$\mathcal{B}(B^- \to D^0 K_2^{*-}) = (2.40^{+1.30+0.72}_{-0.97-1.05}) \times 10^{-5},$$
$$\mathcal{B}(B^- \to \bar{D}^0 K_2^{*-}) = (3.32^{+1.90+1.07}_{-1.18-0.74}) \times 10^{-6},$$

where the first uncertainties are from f_B and ω_b in the B meson wave functions, the second errors are from $\Lambda_{\rm QCD}$ and the scales defined in the Appendix. (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. [8], obtained under the factorization approach. The main reason is due to the enhancement of nonfactorizable contributions in color-suppressed D^0 emission. In Ref. [34], the authors also predict the branching ratios of $B^- \to (D^0, \bar{D}^0) K_2^{*-}$ under the same approach, which are confirmed here. Our results of ratios and phases of the amplitudes are

$$r_{K_0^*} = 0.24^{+0.02+0.07}_{-0.01-0.04}, \qquad \delta_{K_0^*} = (-125.65^{+3.95+23.16}_{-0.00-17.42})^\circ, \quad S1$$
$$r_{K_0^*} = 0.54^{+0.03+0.11}_{-0.02-0.07}, \qquad \delta_{K_0^*} = (-161.51^{+0.91+12.01}_{-0.16-9.53})^\circ, \quad S2$$
$$r_{K_2^*} = 0.37^{+0.02+0.17}_{-0.00-0.09}, \qquad \delta_{K_2^*} = (155.53^{+0.00+2.98}_{-3.36-3.49})^\circ. \tag{20}$$

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

$$R_{CP\pm}^{K_J} = 2\frac{\mathcal{B}(B^- \to D_{CP\pm} K_J^-) + \mathcal{B}(B^+ \to D_{CP\pm} K_J^+)}{\mathcal{B}(B^- \to D^0 K_J^-) + \mathcal{B}(B^+ \to \bar{D}^0 K_J^+)} = 1 + (r_B^{K_J})^2 \pm 2 r_B^{K_J} \cos\delta_B^{K_J} \cos\gamma,$$
$$A_{CP\pm}^{K_J} = \frac{\mathcal{B}(B^- \to D_{CP\pm} K_J^-) - \mathcal{B}(B^+ \to D_{CP\pm} K_J^+)}{\mathcal{B}(B^- \to D_{CP\pm} K_J^-) + \mathcal{B}(B^+ \to D_{CP\pm} K_J^+)} = \pm 2 r_B^{K_J} \sin\delta_B^{K_J} \sin\gamma / R_{CP\pm}^{K}. \tag{21}$$

In the limit of $r_B \to 0$, the ratio $R_{CP\pm}^{K}$ 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 α_s) between the QCD vacuum and the $K_0^*(K_2^*)$ state is small (identically zero), the low sensitivity to γ is improved and in particular large CP asymmetries are expected. The dependence of R_{CP} and A_{CP} on γ is shown in Fig. 2. Since the errors of $r_{K^*_{0,2}}$ and $\delta_{K^*_{0,2}}$ are not large, only their central values are used. We investigate these observables in the region $\gamma = (68^{+10}_{-11})^\circ$ which is from a combined analysis of $B^\pm \to DK^\pm$ [7]. In this region we find that the observables of the $B^- \to (\bar{D}^0, D^0)K_0^{*-}$ 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 γ extracted from the $B^\pm \to DK^\pm$ decays.

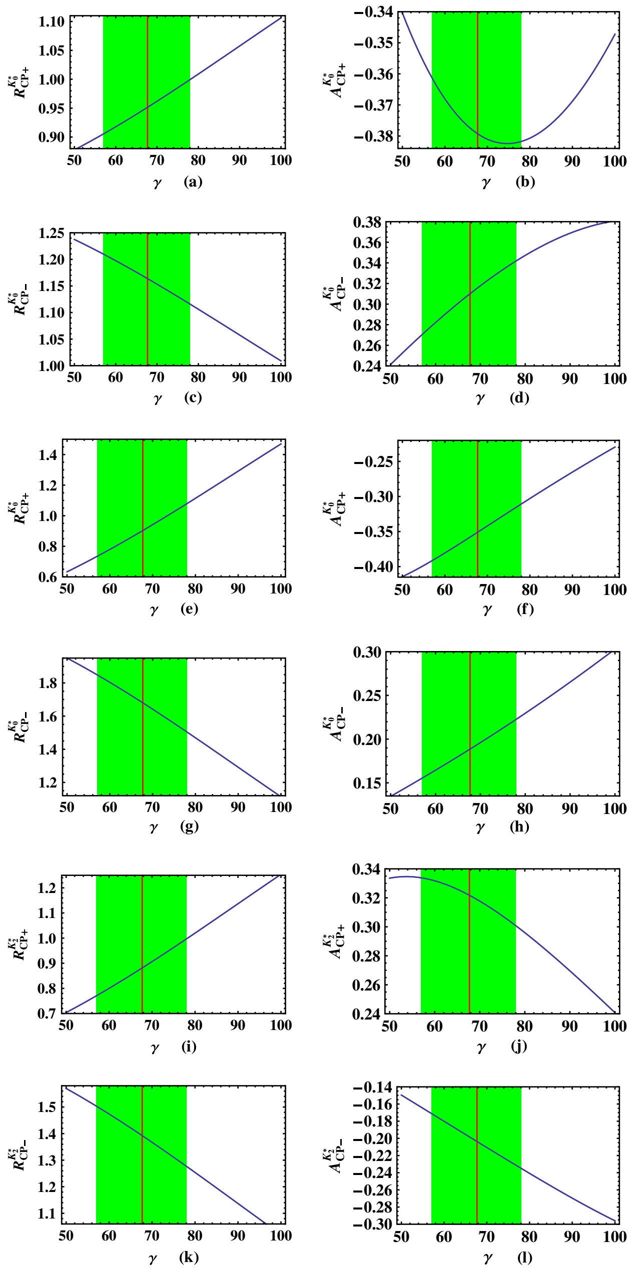

FIG. 2 (color online). The dependence of R_{CP} and A_{CP} on γ. Diagrams (a)–(d) show $R_{CP}^{K_0^*}$ and $A_{CP}^{K_0^*}$ in $S1$, (e)–(h) in $S2$, and diagrams (i)–(l) show $R_{CP}^{K_2^*}$ and $A_{CP}^{K_2^*}$. The shadowed (green) region denotes the current bounds on $\gamma = (68^{+10}_{-11})^\circ$ from a combined analysis of $B^\pm \to DK^\pm$ [7], and the vertical (red) line represents the central value.

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 $\mathcal{O}(A_{CP}^{\mathrm{dir}})/r_{K^*_{0,2}}$ [6] 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 γ 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 Figs. 1(e)–1(h) contribute, in which the spectator quark $\bar{u}$ needs to be replaced by $\bar{s}$ and the c and $\bar{u}$ in the emission meson should be exchanged for the $B_s \to \bar{D}(f_0, f_2')$ decays. The amplitudes are given by

$$\begin{aligned} A(\bar{B}_s^0 \to \bar{D}^0(f_0, f_2')) &= \frac{G_F}{\sqrt{2}} V_{ub} V_{cs}^* (\xi_{\rm in}' + \mathcal{M}_{\rm in}'), \\ A(\bar{B}_s^0 \to D^0(f_0, f_2')) &= \frac{G_F}{\sqrt{2}} V_{cb} V_{us}^* (\xi_{\rm in} + \mathcal{M}_{\rm in}). \end{aligned} \tag{22}$$

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]

$$\begin{aligned} f_{B_s} &= 0.2420 \pm 0.0095, & \bar{f}_{f_0} &= 0.37 \pm 0.02, \\ B_1(f_0) &= -0.78 \pm 0.08, & B_3(f_0) &= 0.02 \pm 0.07, \\ f_{f_2'} &= 0.126 \pm 0.004, & f_{f_2'}^T &= 0.065 \pm 0.012, \end{aligned} \tag{23}$$

with which the branching ratios are predicted as

$$\begin{aligned} \mathcal{B}(\bar{B}_s^0 \to D^0 f_0) &= (3.50^{+1.26+0.56}_{-1.15-0.77}) \times 10^{-5}, \\ \mathcal{B}(\bar{B}_s^0 \to \bar{D}^0 f_0) &= (5.94^{+3.13+1.47}_{-2.16-0.97}) \times 10^{-6}, \\ \mathcal{B}(\bar{B}_s^0 \to D^0 f_2') &= (1.08^{+0.37+0.26}_{-0.35-0.28}) \times 10^{-5}, \\ \mathcal{B}(\bar{B}_s^0 \to \bar{D}^0 f_2') &= (2.85^{+1.53+0.67}_{-0.95-0.43}) \times 10^{-6}. \end{aligned} \tag{24}$$

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 γ.

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 the $B \to DK$ 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.

ACKNOWLEDGMENTS

The work of C. S. K. and R. H. L. is supported by the National Research Foundation of Korea (NRF) grant funded by Korea government of the Ministry of Education, Science and Technology (MEST) (Grants No. 2011-0017430 and No. 2011-0020333). The work of W. W. is supported in part by the DFG and the NSFC through funds provided to the Sino-German CRC 110 "Symmetries and the Emergence of Structure in QCD." The computation code can be found: http://www.itkp.uni-bonn.de/~weiwang/.

APPENDIX: HARD KERNELS IN THE PQCD CALCULATION

The offshellness of the intermediate gluon

$$Q_{a,b,c,d} = x_1\bar{x}_3 m_B^2, \qquad Q_{e,f,g,h,g',h'} = x_1\bar{x}_3(1-r_D^2)m_B^2,$$
$$Q_{k,l,m,n} = x_3\bar{x}_2(1-r_D^2)m_B^2,$$

and the quarks

$$\begin{aligned}
P_a &= \bar{x}_3 m_B^2, \qquad P_b = x_1 m_B^2,\\
P_c &= \bar{x}_3[x_1 - x_2(1-r_D^2)]m_B^2,\\
P_d &= \bar{x}_3[x_1 - \bar{x}_2(1-r_D^2)]m_B^2,\\
P_e &= \bar{x}_3(1-r_D^2)m_B^2, \qquad P_f = x_1(1-r_D^2)m_B^2,\\
P_g &= \{(x_1-x_2)[(1-r_D^2)\bar{x}_3 + r_D^2] + r_D^2\}m_B^2,\\
P_h &= \{(x_1-\bar{x}_2)(1-r_D^2)\bar{x}_3\}m_B^2,\\
P'_g &= \{(x_1-x_2)(1-r_D^2)\bar{x}_3\}m_B^2,\\
P_h &= \{(x_1-\bar{x}_2)[(1-r_D^2)\bar{x}_3 + r_D^2] + r_D^2\}m_B^2,\\
P_k &= x_3(1-r_D^2)m_B^2, \qquad P_l = \bar{x}_2(1-r_D^2)m_B^2,\\
P_m &= -\{\bar{x}_3[1-(1-r_D^2)\bar{x}_2 - x_1] - 1\}m_B^2,\\
P_n &= x_3[x_1 - \bar{x}_2(1-r_D^2)]m_B^2,
\end{aligned}$$

result in the hard scales and kernels

$$\begin{aligned}
t_i &= \max\{\sqrt{Q_i}, \sqrt{P_i}, 1/b_1, 1/b_3\},\\
t_j &= \max\left\{\sqrt{Q_j}, \sqrt{P_j}, 1/b_1, 1/b_2\right\},\\
t_x &= \max\{\sqrt{Q_x}, \sqrt{P_x}, 1/b_2, 1/b_3\},\\
t_y &= \max\left\{\sqrt{Q_y}, \sqrt{P_y}, 1/b_1, 1/b_2\right\},\\
h_{a,e} &= H_e(P_{a,e}, Q_{a,e}, b_1, b_3)S_t(x_3),\\
h_{b,f} &= H_e(P_{b,f}, Q_{b,f}, b_3, b_1)S_t(x_1),\\
h_j &= H_{en}(Q_j, P_j, b_1, b_2),\\
h_k &= H_{af}(P_k, Q_k, b_2, b_3)S_t(x_3),\\
h_l &= H_{af}(P_l, Q_l, b_3, b_2)S_t(x_2),\\
h_{m,n} &= H_{an},
\end{aligned} \tag{A1}$$

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

$$\begin{aligned}
H_e(\alpha,\beta,b_1,b_3) &= K_0(\sqrt{\beta}b_1)[\theta(b_1-b_3)I_0(\sqrt{\alpha}b_3)\\
&\quad\times K_0(\sqrt{\alpha}b_1) + (b_1\leftrightarrow b_3)],\\
H_{en}(\alpha,\beta,b_1,b_2) &= [\theta(b_2-b_1)K_0(\sqrt{\alpha}b_2)\\
&\quad\times I_0(\sqrt{\alpha}b_1) + (b_1\leftrightarrow b_2)]\\
&\quad\times\begin{cases}\frac{i\pi}{2}H_0^{(1)}(\sqrt{\beta}b_2), & \beta<0\\ K_0(\sqrt{\beta}b_2), & \beta>0,\end{cases}
\end{aligned} \tag{A2}$$

$$\begin{aligned}
H_{af}(\alpha,\beta,b_2,b_3) &= \left(i\frac{\pi}{2}\right)^2 H_0^{(1)}(\sqrt{\beta}b_2)[\theta(b_2-b_3)\\
&\quad\times H_0^{(1)}(\sqrt{\alpha}b_2)J_0(\sqrt{\alpha}b_3)\\
&\quad+\theta(b_3-b_2)H_0^{(1)}(\sqrt{\alpha}b_3)J_0(\sqrt{\alpha}b_2)],\\
H_{an} &= i\frac{\pi}{2}\Big[\theta(b_1-b_2)H_0^{(1)}\Big(\sqrt{Q_y}b_1\Big)J_0\Big(\sqrt{Q_y}b_2\Big)\\
&\quad+\theta(b_2-b_1)H_0^{(1)}\Big(\sqrt{Q_y}b_2\Big)J_0\Big(\sqrt{Q_y}b_1\Big)\Big]\\
&\quad\times\begin{Bmatrix} K_0\Big(\sqrt{|P_y|}b_1\Big) & \text{for } P_y\le 0\\ \frac{i\pi}{2}H_0^{(1)}\Big(\sqrt{|P_y|}b_1\Big) & \text{for } P_y\ge 0\end{Bmatrix},
\end{aligned} \tag{A3}$$

where $H_0^{(1)}(z) = \mathrm{J}_0(z) + i\mathrm{Y}_0(z)$. The S_t has been parameterized as

$$S_t(x) = \frac{2^{1+2c}\Gamma(3/2+c)}{\sqrt{\pi}\Gamma(1+c)}[x(1-x)]^c, \tag{A4}$$

with $c = 0.4 \pm 0.1$.

The E_is contain the Sudakov evolution factors and the strong coupling constants, and are given by

$$E_{a,b}(t) = \alpha_s(t)\exp[-S_B(t) - S_D(t, \bar{x}_3, b_3)],$$
$$E_{c,d}(t) = \alpha_s(t)\exp[-S_B(t) - S_D(t, \bar{x}_3, b_1) - S_{K_0^*}(t, b_2)],$$
$$E_{e,f}(t) = \alpha_s(t)\exp[-S_B(t) - S_{K_0^*}(t, b_3)],$$
$$E_{g,h}(t) = \alpha_s(t)\exp[-S_B(t) - S_D(t, \bar{x}_2, b_2) - S_{K_0^*}(t, b_1)],$$
$$E_{g',h'}(t) = \alpha_s(t)\exp[-S_B(t) - S_D(t, x_2, b_2) - S_{K_0^*}(t, b_1)],$$
$$E_{k,l}(t) = \alpha_s(t)\exp[-S_D(t, x_3, b_3) - S_{K_0^*}(t, b_2)],$$
$$E_{m,n}(t) = \alpha_s(t)\exp[-S_B(t) - S_D(t, x_2, b_2) - S_{K_0^*}(t, b_2)].$$

in which the Sudakov exponents are defined as

$$S_B(t) = s\left(x_1\frac{m_B}{\sqrt{2}}, b_1\right) + \frac{5}{3}\int_{1/b_1}^{t}\frac{d\bar{\mu}}{\bar{\mu}}\gamma_q(\alpha_s(\bar{\mu})),$$
$$S_D(t, x_2, b_2) = s\left(x_2\frac{m_B}{\sqrt{2}}, b_2\right) + 2\int_{1/b_2}^{t}\frac{d\bar{\mu}}{\bar{\mu}}\gamma_q(\alpha_s(\bar{\mu})),$$
$$S_{K_0^*}(t, b_3) = s\left(x_3\frac{m_B}{\sqrt{2}}, b_3\right) + s\left((1-x_3)\frac{m_B}{\sqrt{2}}, b_3\right) + 2\int_{1/b_3}^{t}\frac{d\bar{\mu}}{\bar{\mu}}\gamma_q(\alpha_s(\bar{\mu})), \tag{A5}$$

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

[1] D. Asner *et al.* (HFAG Collaboration), arXiv:1207.1158; updated results are available at http://www.slac.stanford.edu/xorg/hfag/.
[2] For recent measurements, please see J.P. Lees *et al.* (*BABAR* Collaboration), Phys. Rev. D **87**, 052015 (2013); M. John, T.M. Karbach *et al.* LHCb Collaboration, Report No. LHCb-CONF-2012-032; K. Trabelsi *et al.* (Belle Collaboration), arXiv:1301.2033; S. Ricciardi, M. Ishino, K. Nagano, S. Asai *et al.* (LHCb Collaboration), EPJ Web Conf. **49**, 13006 (2013).
[3] M. Gronau and D. Wyler, Phys. Lett. B **265**, 172 (1991).
[4] M. Gronau and D. London, Phys. Lett. B **253**, 483 (1991).
[5] I. Dunietz, Phys. Lett. B **270**, 75 (1991).
[6] W. Wang, Phys. Rev. Lett. **110**, 061802 (2013); M. Martone and J. Zupan, Phys. Rev. D **87**, 034005 (2013); B. Bhattacharya, D. London, M. Gronau, and J.L. Rosner, Phys. Rev. D **87**, 074002 (2013).
[7] J. Charles, A. Höcker, H. Lacker, S. Laplace, F.R. Diberder, J. Malclés, J. Ocariz, M. Pivk, and L. Roos (CKMfitter Group), Eur. Phys. J. C **41**, 1 (2005); updated results and plots are available at http://ckmfitter.in2p3.fr.
[8] W. Wang, Phys. Rev. D **85**, 051301 (2012); AIP Conf. Proc. **1492**, 117 (2012).
[9] M. Diehl and G. Hiller, J. High Energy Phys. 06 (2001) 067; Phys. Lett. B **517**, 125 (2001).
[10] H.-n. Li and H.-L. Yu, Phys. Rev. D **53**, 2480 (1996).
[11] Y. Y. Keum, H.-N. Li, and A. I. Sanda, Phys. Lett. B **504**, 6 (2001); Phys. Rev. D **63**, 054008 (2001).
[12] C.-D. Lu, K. Ukai, and M.-Z. Yang, Phys. Rev. D **63**, 074009 (2001).
[13] H.-n. Li, Y.-L. Shen, and Y.-M. Wang, Phys. Rev. D **85**, 074004 (2012); J. High Energy Phys. 02 (2013) 008.
[14] W.-F. Wang and Z.-J. Xiao, Phys. Rev. D **86**, 114025 (2012); Y.-Y. Fan, W.-F. Wang, and Z.-J. Xiao, arXiv:1301.6246.
[15] T. Kurimoto, H.-n. Li, and A. I. Sanda, Phys. Rev. D **67**, 054028 (2003).
[16] Y.-Y. Keum, T. Kurimoto, H. N. Li, C.-D. Lu, and A. I. Sanda, Phys. Rev. D **69**, 094018 (2004).
[17] C.-D. Lu and K. Ukai, Eur. Phys. J. C **28**, 305 (2003).
[18] Y. Li and C.-D. Lu, J. Phys. G **29**, 2115 (2003).
[19] C.-D. Lu, Phys. Rev. D **68**, 097502 (2003).
[20] R.-H. Li, C.-D. Lu, and H. Zou, Phys. Rev. D **78**, 014018 (2008).
[21] H. Zou, R.-H. Li, X.-X. Wang, and C.-D. Lu, J. Phys. G **37**, 015002 (2010).
[22] R.-H. Li, X.-X. Wang, A. I. Sanda, and C.-D. Lu, Phys. Rev. D **81**, 034006 (2010).
[23] H.-n. Li, Prog. Part. Nucl. Phys. **51**, 85 (2003).
[24] C.-D. Lu and M.-Z. Yang, Eur. Phys. J. C **28**, 515 (2003).
[25] H.-Y. Cheng, C.-K. Chua, and K.-C. Yang, Phys. Rev. D **73**, 014017 (2006); H.-Y. Cheng, C.-K. Chua, K.-C. Yang, and Z.-Q. Zhang, Phys. Rev. D **87**, 114001 (2013).
[26] C.-D. Lu, Y.-M. Wang, and H. Zou, Phys. Rev. D **75**, 056001 (2007).
[27] H.-Y. Han, X.-G. Wu, H.-B. Fu, Q.-L. Zhang, and T. Zhong, Eur. Phys. J. A **49**, 78 (2013).
[28] R. H. Li, C. D. Lu, W. Wang, and X. X. Wang, Phys. Rev. D **79**, 014013 (2009).
[29] H.-Y. Cheng, Y. Koike, and K.-C. Yang, Phys. Rev. D **82**, 054019 (2010).
[30] E. T. Neil *et al.* (Fermilab Lattice and MILC Collaborations), Proc. Sci., LATTICE2011 (2011) 320 [arXiv:1112.3978].
[31] J. Beringer *et al.* (Particle Data Group Collaboration), Phys. Rev. D **86**, 010001 (2012).
[32] W. Wang, Y.-L. Shen, Y. Li, and C.-D. Lu, Phys. Rev. D **74**, 114010 (2006); Y.-L. Shen, W. Wang, J. Zhu, and C.-D. Lu, Eur. Phys. J. C **50**, 877 (2007); X. Liu, Z.-Q. Zhang, and Z.-J. Xiao, Chinese Phys. C **34**, 157 (2010); C. S. Kim, Y. Li, and W. Wang, Phys. Rev. D **81**, 074014 (2010); X. Liu and Z.-J. Xiao, Commun. Theor. Phys. **53**, 540 (2010); X. Liu, Z.-J. Xiao, and Z.-T. Zou, arXiv:1105.5761; Z.-Q. Zhang, Phys. Rev. D **82**, 034036 (2010).
[33] W. Wang, Phys. Rev. D **83**, 014008 (2011).
[34] Z.-T. Zou, X. Yu, and C.-D. Lu, Phys. Rev. D **86**, 094015 (2012); **86**, 094001 (2012).
[35] S. Freddy, C. S. Kim, R. H. Li, and Z. T. Zou, Phys. Rev. D **88**, 014031 (2013).