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Decay modes $\overline{B^0_{d,s}} \to \gamma D^{*0}$

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The various decay modes of the type $B \to \gamma D^*$ are dynamically different. In general, there are factorized contributions of pole and nonpole type, and pseudoscalar exchange contributions at meson level. The purpose of this paper is to point out that the decay modes $\overline{B^0_{d,s}} \to \gamma D^{*0}$ have small contributions from such mechanisms, in contrast to the decay modes $\overline{B^0_{d,s}} \to \gamma \overline{D^{*0}}$ and $B^- \to \gamma D^{*-}_{s,d}$. On the other hand, there are nonfactorizable $1/N_c$ suppressed contributions from colored four quark operators obtained from Fierz transformations of the standard four quark operators. Such $1/N_c$ suppressed terms involving emission of soft gluons are calculated within a heavy-light chiral quark model, and they are found to dominate the amplitudes for the decay modes $\overline{B^0_{d,s}} \to \gamma D^{*0}$. We estimate the branching ratio for these modes in the low velocity regime, i.e. in the heavy quark limit, both for the b and the c quarks. In this limit we obtain a value $\simeq 1 \times 10^{-5}$ for $\overline{B^0_d} \to \gamma D^{*0}$, and $\simeq 6 \times 10^{-7}$ for $\overline{B^0_s} \to \gamma D^{*0}$. We expect substantial corrections to these numbers because the energy gap between the b- and c-quark masses are significantly bigger than 1 GeV. However, we expect that our estimates of the *amplitudes* for $\overline{B^0_{d,s}} \to \gamma D^{*0}$ give the right order of magnitude because other mechanisms are numerically suppressed in this special case.

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I. INTRODUCTION

There is presently great interest in decays of B-mesons, due to numerous experimental results coming from *BABAR* and Belle. Later LHC will provide data for such processes. B-decays of the type $B \to \pi\pi$ and $B \to K\pi$, where the energy release is big compared to the light meson masses, has been treated within *QCD factorization* and *soft collinear effective theory* (SCET) [1]. In these cases the amplitudes factorize into products of two matrix elements of weak currents in the high energy limit, and nonfactorizable corrections of order α_s can be calculated perturbatively.

The decays $B \to \pi\pi$, $K\pi$ are typical heavy to light decays. It was pointed out in previous papers [2] that for various decays of the type $\bar{B} \to D\bar{D}$, which are of heavy to heavy type, the methods of [1] are not expected to hold because the energy release is of order 1 GeV only. (Here $\bar{B}$, D, and $\bar{D}$ contain a heavy b, c, and anti-c quark, respectively.) In this paper we consider decay modes of the type $B \to \gamma D^*$. Such modes have been studied in the literature [3–5] for some time. We restrict ourselves to processes where the b-quark decays. This means the quark level processes $b\bar{q} \to \gamma c\bar{u}$, $b\bar{q} \to \gamma u\bar{c}$, and $b\bar{u} \to \gamma c\bar{q}$, where $q = d$ or $q = s$. Processes where the anti-b-quark decays proceed analogously.

Formally, decays of the type $B \to \gamma D$ are heavy to heavy transitions in the heavy quark limits $(1/m_b) \to 0$ *and* $(1/m_c) \to 0$, and in Refs. [4,5] the decay of a charged B-meson was studied within *heavy quark effective theory* (HQEFT) [6] and *heavy-light chiral perturbation theory* (HLχPT) [7]. This framework was also used to study the Isgur-Wise function for the $B \to D$ transition currents, which is also a heavy to heavy transition where chiral loops (in terms of HLχPT) and $1/m_{b,c}$ corrections (in terms of HQEFT) have been added [8]. In the present paper, we will also stick to this framework, although it is not expected to hold for precise numerical estimates because the energy gap between the b- and the c-scale is substantial, namely, about 3 times the chiral symmetry breaking scale. However, for $B \to \gamma D$ there is no ideal framework. The other extreme would be to consider the limit where the c-quark is light as the u, d, s-quarks. Still our framework can be used as a classification scheme. And we will use it for estimating $\overline{B^0_{d,s}} \to \gamma D^{*0}$ where only the soft gluon mechanism described in Sec. III is expected to give the dominant contribution.

Classifying decay modes of the type $B \to \gamma D^*$, they might have substantial factorized contributions, of pole or nonpole type. Second, there are also meson exchange contributions. These will be chiral loop contributions in the HQEFT limit (for both b and c-quarks). Such meson exchange diagrams, which are nonfactorizable and $1/N_c$ suppressed, are significant for the decay modes $\overline{B^0_{d,s}} \to \gamma \overline{D^{*0}}$ and $B^- \to \gamma D^{*-}_{s,d}$. They could be handled with dispersion relations if the decays $B \to DM$ ($M =$ light meson) were known in more detail.

The purpose of this paper is first to point out that the decay modes $\overline{B^0_{d,s}} \to \gamma D^{*0}$ have almost zero contributions from factorized and meson exchange amplitudes. Second, these decay modes have significant contributions from soft gluon emission [9]. Such nonfactorizable (color suppressed $\sim 1/N_c$) contributions to $B - \bar{B}$ mixing [10], $B \to D\bar{D}$ [2], and $B \to D\eta'$ [11] decays have been calculated in terms of the (lowest dimension, model dependent) gluon condensate within a recently developed *heavy-light chiral*

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quark model (HLχQM) [12], which is based on the HQEFT [6] and HLχPT [7]. At this point we extend the framework of [4,5] in order to estimate the coefficients of $1/N_c$ suppressed terms within HLχPT. Soft gluon emission has also been considered within the chiral quark model in the pure light sector to estimate $K-\bar{K}$ mixing and $K\to 2\pi$ decays [13]. We estimate the branching ratios for $\overline{B^0_{d,s}}\to\gamma D^{*0}$ in the heavy b- and c-quark limits. Note that the decay modes $\overline{B^0_{d,s}}\to\gamma D^{*0}$ and $\overline{B^0_{d,s}}\to\gamma\overline{D^{*0}}$ proceed differently. In the last case there are significant meson exchange contributions.

In Sec. II we present the weak four quark Lagrangian and its factorized and nonfactorizable matrix elements. In Sec. III we present the framework of HQEFT, HLχPT, and HLχQM. In Sec. IV we calculate explicitly factorizable contributions, and more important for the purpose of this letter, the nonfactorizable matrix elements due to soft gluons expressed through the (model dependent) quark condensate. In Sec. V we give the results and conclusion.

II. THE WEAK QUARK LAGRANGIAN AND ITS MATRIX ELEMENTS

Based on the electroweak and quantum chromodynamical interactions, one constructs an effective nonleptonic Lagrangian at quark level in the standard way:

$$\mathcal{L}_W=\sum_i C_i(\mu)Q_i(\mu), \tag{1}$$

where all information of the short distance (SD) loop effects above a renormalization scale μ is contained in the Wilson coefficients C_i. In our case there are four relevant operators,

$$\begin{aligned}Q_1&=4(\bar{q}_L\gamma^\alpha b_L)(\bar{c}_L\gamma_\alpha u_L),\\ Q_2&=4(\bar{c}_L\gamma^\alpha b_L)(\bar{q}_L\gamma_\alpha u_L),\end{aligned} \tag{2}$$

$$\begin{aligned}Q_3&=4(\bar{q}_L\gamma^\alpha b_L)(\bar{u}_L\gamma_\alpha c_L),\\ Q_4&=4(\bar{u}_L\gamma^\alpha b_L)(\bar{q}_L\gamma_\alpha c_L),\end{aligned} \tag{3}$$

for $q=d$, s. This effective Lagrangian is based on the interactions in Fig. 1 and hard gluon corrections to these diagrams. Operators obtained from penguin diagrams have

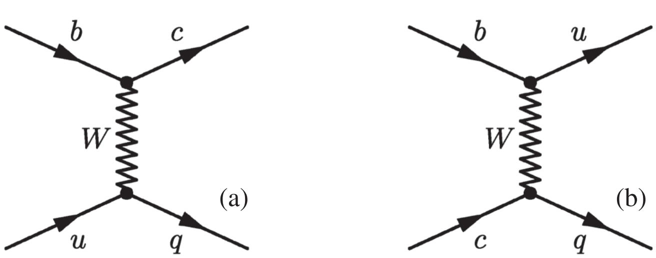

FIG. 1. Tree level W-exchange leading to the effective Lagrangian in Eq. (1). The left diagram (a) gives rise to $Q_{1,2}$, and the right diagram (b) gives rise to $Q_{3,4}$.

small Wilson coefficients and will be omitted in the present paper. The coefficients $C_{1,2}$ and $C_{3,4}$ have different Kobayashi-Maskawa (KM) quark mixing structures. We may write

$$C_i=-\frac{G_F}{\sqrt{2}}(V_{cb}V^*_{uq})a_i;\qquad C_j=-\frac{G_F}{\sqrt{2}}(V_{ub}V^*_{cq})a_j, \tag{4}$$

for $i=1,2$ and $j=3,4$, respectively. Here the "reduced" Wilson coefficients a_i (for $i=1,2,3,4$) are dimensionless numbers. (In practice, $a_1=a_3$ and $a_2=a_4$.) Furthermore, in terms of the Wolfenstein parameter λ, we have $V_{cb}V^*_{ud}\sim\mathcal{O}(\lambda^2)$. For $q=s$ the KM factors $V_{cb}V^*_{us}$ and $V_{ub}V^*_{cs}$ are both $\sim\mathcal{O}(\lambda^3)$, while $V_{ub}V^*_{cd}\sim\mathcal{O}(\lambda^4)$. At the scale $\mu=M_W$, when perturbative QCD is switched off, one has $a_{1,3}=0$ and $a_{2,4}=1$. At the scale $\mu=m_b$, $a_{1,3}\sim 10^{-1}$ and negative, and $a_{2,4}$ are slightly bigger than one [14]. Extrapolating the Wilson coefficients down to $\mu\sim\Lambda_\chi\sim 1$ GeV, which is the matching scale between short and long distance effects within our framework [2,10,11], we obtain $a_{2,4}\simeq 1.17$ and $a_{1,3}\simeq -0.37$ [11]. Alternatively, one might perform perturbative QCD corrections within HQEFT as done in [15] and used in [2] for $B\to D\bar{D}$, but numerical differences will be small.

For $B\to D\gamma$, one may also think of operators like

$$\begin{aligned}&eF_{\mu\nu}(\bar{q}_L\gamma^\mu b_L)(\bar{c}_L\gamma^\nu u_L),\\ &e(\bar{q}_L\sigma^{\alpha\beta}F_{\alpha\beta}\gamma^\mu b_L)(\bar{c}_L\gamma_\mu u_L).\end{aligned} \tag{5}$$

However, such operators are of dimension eight, and dominated at low momenta which make a short distance treatment dubious.

In the *factorized* limit, where there are no strong interactions between the two quark currents in Q_i, we obtain the amplitude for $\overline{B^0_q}\to\gamma D^{*0}$ from (1)–(3):

$$\begin{aligned}\langle\gamma D^{*0}|\mathcal{L}_W|\overline{B^0_q}\rangle_F=4\Big(C_1+\frac{C_2}{N_c}\Big)&(\langle D^{*0}|\overline{c_L}\gamma_\mu u_L|0\rangle\\ &\times\langle\gamma|\overline{q_L}\gamma_\mu b_L|\overline{B^0}\rangle\\ &+\langle\gamma D^{*0}|\overline{c_L}\gamma_\mu u_L|0\rangle\langle 0|\overline{q_L}\gamma_\mu b_L|\overline{B^0}\rangle),\end{aligned} \tag{6}$$

where the subscript F means "factorized." For $\overline{B^0_q}\to\gamma\overline{D^{*0}}$ we obtain the same expression with $C_{1,2}$ replaced by $C_{3,4}$ and with c and u interchanged. Thus, the neutral decays have *small factorized contributions* proportional to $a_{nf}=(a_{1,3}+a_{2,4}/N_c)$, which is $\sim 10^{-2}$.

For the charged case $B^- \to \gamma D_q^{*-}$ we obtain

$$\begin{aligned}\langle \gamma D_q^{*-}|\mathcal{L}_W|B^-\rangle_F = 4\left(C_4 + \frac{C_3}{N_c}\right)&(\langle D_q^{*-}|\overline{q_L}\gamma_\mu c_L|0\rangle \\ &\times \langle \gamma|\overline{u_L}\gamma_\mu b_L|B^-\rangle \\ &+ \langle D_q^{*-}\gamma|\overline{q_L}\gamma_\mu c_L|0\rangle \\ &\times \langle 0|\overline{u_L}\gamma_\mu b_L|B^-\rangle).\end{aligned} \quad (7)$$

The charged decays have substantial factorized contributions proportional to $a_f = (a_4 + a_3/N_c) \sim 1$. These are visualized in Fig. 2.

In order to study nonfactorizable contributions at quark level, we use the following relation between the generators of $SU(3)_c$ (i, j, l, n are color indices running from 1 to 3):

$$\delta_{ij}\delta_{ln} = \frac{1}{N_c}\delta_{\text{in}}\delta_{lj} + 2t^a_{\text{in}}t^a_{lj}, \quad (8)$$

where a is the color octet index. Then the operators $Q_{1,2}$ may, by means of a Fierz transformation, be written in the following way:

$$Q_{1,3} = \frac{1}{N_c}Q_{2,4} + 2\tilde{Q}_{2,4}, \qquad Q_{2,4} = \frac{1}{N_c}Q_{1,3} + 2\tilde{Q}_{1,3}, \quad (9)$$

where the operators with the "tilde" contain color matrices:

$$\begin{aligned}\tilde{Q}_1 &= 4(\bar{q}_L\gamma^\alpha t^a b_L)(\bar{c}_L\gamma_\alpha t^a u_L), \\ \tilde{Q}_2 &= 4(\bar{c}_L\gamma^\alpha t^a b_L)(\bar{q}_L\gamma_\alpha t^a u_L).\end{aligned} \quad (10)$$

$$\begin{aligned}\tilde{Q}_3 &= 4(\bar{q}_L\gamma^\alpha t^a b_L)(\bar{u}_L\gamma_\alpha t^a c_L), \\ \tilde{Q}_4 &= 4(\bar{u}_L\gamma^\alpha t^a b_L)(\bar{q}_L\gamma_\alpha t^a c_L).\end{aligned} \quad (11)$$

Note that the C_i/N_c terms in (6) and (7) stem from the $1/N_c$ terms in (9).

The nonfactorizable amplitude for $\overline{B^0_q} \to \gamma D^{*0}$, with one gluon emission obtained from the colored operators in (10) and (11), might be written in a quasifactorizable way in terms of octet gluonic intermediate states ($|G\rangle$):

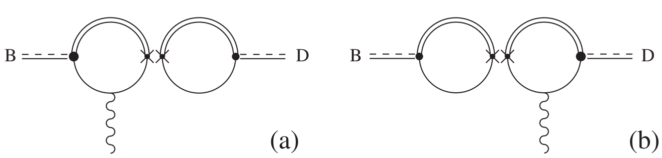

FIG. 2. Factorized contributions for $B \to \gamma D^*$. The combined dashed and full lines represent heavy mesons, the double lines represent heavy quarks, and the single lines light quarks. The wavy line is a photon. (a) Emission of a photon from the B-meson. (b) Emission of a photon from the D-meson.

$$\begin{aligned}\langle \gamma D^{*0}|\mathcal{L}_W|\overline{B^0_q}\rangle_{\text{NFG}} &= 8C_2\langle \gamma D^{*0}|\tilde{Q}_1|\overline{B^0_q}\rangle \\ &= 8C_2(\langle D^{*0}|\overline{c_L}\gamma_\mu t^a u_L|G\rangle \\ &\quad \times \langle G\gamma|\overline{q_L}\gamma_\mu t^a b_L|\overline{B^0_q}\rangle \\ &\quad + \langle \gamma D^{*0}|\overline{c_L}\gamma_\mu t^a u_L|G\rangle \\ &\quad \times \langle G|\overline{q_L}\gamma_\mu t^a b_L|\overline{B^0_q}\rangle).\end{aligned} \quad (12)$$

This amplitude is visualized in Fig. 3. It has to be calculated within some framework describing long distance gluonic effects. In our case we have chosen the HLχQM [12].

The nonfactorizable amplitude for $\overline{B^0_q} \to \gamma \overline{D^{*0}}$ with one gluon emission obtained from the colored operators is the same as above with $C_2 \to C_4$ and with u and c-quarks interchanged. The nonfactorizable amplitude for $B^- \to \gamma D_q^{*-}$ with gluon emission from the colored operator $\tilde{Q}_4$ is proportional to a_3 and therefore relatively small.

We observe the following generic pattern: Some decay modes have substantial factorizable contributions proportional to the favorable Wilson coefficient linear combination $a_f \equiv (a_{2,4} + a_{1,3}/N_c)$, which is close to 1. In this case there are contributions from the colored operators $\tilde{Q}_{1,3}$ proportional to $a_{1,3}$ of minor importance. For other modes there might be factorized matrix elements proportional to the nonfavorable coefficient $a_{nf} \equiv (a_{1,3} + a_{2,4}/N_c)$ which is close to zero (of order 10^{-2} or smaller) at our matching scale $\mu = \Lambda_\chi$. In these cases there are significant contributions proportional to $a_{2,4} \sim 1$ from the operators $\tilde{Q}_{1,3}$, involving color exchange.

In terms of the B-meson pseudoscalar field Φ, the D^*-meson vector field V^μ, and the electromagnetic field tensor $F_{\mu\nu}$, we can write down the effective Lagrangian to first order in the photon momentum, consistent with the heavy quark limits:

$$\mathcal{L}_{\text{eff}} = A^{(+)}i\epsilon_{\mu\nu\alpha\beta}\Phi F^{\mu\nu}V^\alpha v_b^\beta + A^{(-)}\Phi F_{\mu\nu}V^\mu v_b^\nu, \quad (13)$$

where the positive and negative parity amplitudes $A^{(\pm)}$ depend on hadronic parameters, and the meson masses $M_{B,D}$.

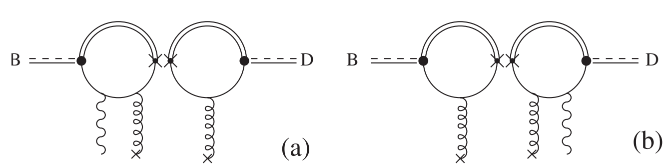

FIG. 3. Nonfactorizable contributions to $B \to \gamma D$ from the colored operators $\tilde{Q}_i$ within a quasifactorized approximation. The curly lines represent soft gluon emission ending in vacuum to make a gluon condensate. (a) With additional photon emission from the B-meson. (b) With additional photon emission from the D-meson.

III. HEAVY-LIGHT CHIRAL LAGRANGIANS FOR $B \to D$ TRANSITIONS

Our calculations will be based on HQEFT [6], which is a systematic $1/m_Q$ expansion in the heavy quark mass m_Q. Each of the heavy quark fields $Q(= b, c\bar{c})$ are replaced with a reduced field $Q_v^{(+)}$ for a heavy quark (b or c), and $Q_v^{(-)}$ for a heavy antiquark (in the case of $\bar{c}$). The Lagrangian for heavy quarks is

$$\mathcal{L}_{\mathrm{HQEFT}} = \pm \overline{Q_v^{(\pm)}} iv \cdot D Q_v^{(\pm)} + \mathcal{O}(m_Q^{-1}), \qquad (14)$$

where v is the velocity of the heavy quark, and D_μ is the covariant derivative containing the gluon and the photon fields. In [10], the $1/m_Q$ corrections were calculated for $B-\bar{B}$-mixing. In this paper such corrections will not be considered.

Integrating out the heavy and light quarks, the effective Lagrangian (the HLχPT terms) up to $\mathcal{O}(m_Q^{-1})$ can be written as [7,12]

$$\mathcal{L} = \mp \mathrm{Tr}[\overline{H_a^{(\pm)}} iv \cdot D_{ba} H_b^{(\pm)}] - g_{\mathcal{A}} \mathrm{Tr}[\overline{H_a^{(\pm)}} H_b^{(\pm)} \gamma_\mu \gamma_5 \mathcal{A}^\mu_{ba}] + \cdots, \qquad (15)$$

where the ellipses denote terms not relevant in this paper. The indices $a, b = 1, 2, 3$ correspond to the quark flavors u, d, s, and $H_a^{(\pm)}$ is the heavy meson field containing a spin zero and spin one boson, and $\mathcal{A}^\mu$ is an axial field:

$$H_a^{(\pm)} \equiv P_\pm (P_{a\mu}^{(\pm)} \gamma^\mu - i P_{a5}^{(\pm)} \gamma_5); \qquad \mathcal{A}_\mu \equiv -\frac{i}{2}(\xi^\dagger \partial_\mu \xi - \xi \partial_\mu \xi^\dagger), \qquad (16)$$

where $P_\pm$ are projecting operators $P_\pm = (1 \pm \gamma \cdot v)/2$, and v is the velocity of the heavy quark. Moreover, $\xi \equiv \exp(i\Pi/f)$, where f is the bare pion coupling, and Π is a 3 by 3 matrix which contains the Goldstone bosons π, K, η in the standard way. The axial chiral coupling is $g_{\mathcal{A}} \simeq 0.6$. Equations (15) and (16) are used for the chiral loop contributions within HLχPT. The covariant derivative is given by $iD^\mu_{ba} = i\delta_{ba}\partial_\mu - e(\tilde{Q}_\xi)_{ba} A^\mu$, where $\tilde{Q}_\xi = \xi Q \xi^\dagger R + \xi^\dagger Q \xi L$, A^μ is the photon field, and $Q = \mathrm{diag}(-2/3, -1/3, -1/3)$ is the electric charge matrix. Further, L and R are the left- and right-handed projection matrices acting in Dirac space.

The relevant Lagrangian term for electromagnetic transition between $B(0^-)$ and $B(1^-)$ mesons [and similarly between $D(0^-)$ and $D(1^-)$ mesons] is given by [5,16,17]

$$\mathcal{L}_\beta = \frac{e\beta}{4} \mathrm{Tr}[\bar{H} H \sigma \cdot F \tilde{Q}_\xi], \qquad (17)$$

where F is the electromagnetic tensor. The constant β is due to radiation from the light quark in the heavy meson. It is not determined within HLχQM alone, but within various quark models it is of order $1/m$, where m is the constituent light quark mass. There is also a similar term $\sim 1/m_Q$ for radiation from the heavy quark in the heavy meson.

The simplest way to calculate the matrix element of four quark operators like Q_{1-4} in Eq. (1) is by inserting vacuum states between the two involved currents, as shown in Eqs. (6) and (7). This is the factorized limit. The matrix elements of the weak currents correspond to the bosonized version of these currents. Based on the symmetry of HQEFT, the bosonized current for decay of the $b\bar{q}$ system is, to lowest order in the chiral expansion [7,12],

$$\overline{q_L} \gamma^\mu Q_{bv}^{(+)} \to \frac{\alpha_H}{2} \mathrm{Tr}[\xi^\dagger \gamma^\alpha L H_{bq}^{(+)}], \qquad (18)$$

where $Q_{bv}^{(+)}$ is a heavy b-quark field, $v = v_b$ is its velocity, and $H_{bq}^{(+)}$ is the corresponding heavy meson field for $\bar{B}_q$. This bosonization has to be compared with the matrix elements defining the meson decay constants $f_H (H = B, D)$. Before perturbative QCD corrections for scales $\mu < m_Q$ and chiral loop corrections have been considered, one has $\alpha_H = f_H \sqrt{M_H}$ (see [6,12]). For the W-boson materializing to a D or $\bar{D}$ mesons, we obtain the bosonized current similar to (18):

$$\overline{Q_{cv}^{(+)}} \gamma^\alpha q_L \to \frac{\alpha_H}{2} \mathrm{Tr}[\gamma^\alpha L \xi \overline{H_{cq}^{(+)}}] \qquad \overline{q_L} \gamma^\alpha Q_{cv}^{(-)} \to \frac{\alpha_H}{2} \mathrm{Tr}[\xi^\dagger \gamma^\alpha L H_{cq}^{(-)}], \qquad (19)$$

where v is the velocity of the heavy c or $\bar{c}$ quarks ($v = v_c$ or $v = v_{\bar{c}}$), and $H_{cq}^{(\pm)}$ is the corresponding field for the D_q or the $\bar{D}_q$ meson.

Up to KM-factors and Wilson coefficients (including Fermi's coupling), the chiral Lagrangian piece for the virtual $B \to D$ transition is now given by the product of the weak currents in (18) and (19). This Lagrangian combined with the term in Eq. (17) is the basis for the pole model calculations in [4,5], where the B or D-meson radiates a photon. Further, there are also direct terms involving the electromagnetic field, for instance occurring as $\sigma_{\mu\nu} F^{\mu\nu}$ inside one of the traces in (18) or (19). (Here $\sigma_{\mu\nu} = i[\gamma_\mu, \gamma_\nu]/2$ as usual.) Also color suppressed terms may occur. Many terms allowed by the symmetry can be written down, but their coefficients are unknown. One way to generate such terms, and to estimate their coefficients, is by means of the HLχQM recently developed in [12], which we will describe shortly in the following.

The HLχQM is especially useful for calculations of matrix elements and chiral Lagrangians for the quark operators in (1) beyond the factorized limit. For this case, this model incorporates emission of soft gluons modeled by a gluon condensate. See also [18,19]. The Lagrangian for the HLχQM is

$$\mathcal{L}_{\mathrm{HL}\chi\mathrm{QM}} = \mathcal{L}_{\mathrm{HQEFT}} + \mathcal{L}_{\chi\mathrm{QM}} + \mathcal{L}_{\mathrm{Int}}. \qquad (20)$$

The first term is given in Eq. (14). The light quark sector is described by the chiral quark model (χQM), having a standard QCD term and a term describing interactions between quarks and (Goldstone) mesons:

$$\mathcal{L}_{\chi\mathrm{QM}} = \bar{\chi}[\gamma^\mu(iD_\mu + \gamma_5 \mathcal{A}_\mu) - m]\chi + \cdots, \qquad (21)$$

where the ellipses denote terms which are irrelevant here. Here m is the $SU(3)$ invariant constituent light quark mass of order 2–300 MeV, and χ is the flavor rotated triplet quark field given by $\chi_L = \xi^\dagger q_L$, $\chi_R = \xi q_R$, where the standard triplet light quark field is given by $q^T = (u, d, s)$. The left- and right-handed projections q_L and q_R are transforming after $SU(3)_L$ and $SU(3)_R$, respectively. The covariant derivative D_μ in (21) contains both the photon field and the soft gluon field forming gluon condensates. The gluon condensate contributions are calculated by Feynman diagram techniques as in [10–13,20,21].

The bosonization (binding) of a heavy quark (heavy antiquark) and light antiquark (light quark) is performed by means of the following interaction Lagrangian [12,19]:

$$\mathcal{L}_{\mathrm{Int}} = -G_H[\bar{\chi}_a \overline{H_a^{(\pm)}} Q_v^{(\pm)} + \overline{Q_v^{(\pm)}} H_a^{(\pm)} \chi_a]. \qquad (22)$$

Here G_H is a coupling constant satisfying $G_H^2 = 2m\rho/f_\pi^2$, where ρ is a hadronic parameter of order one. In [12] it was shown how (15) could be obtained from the HLχQM. Performing this bosonization of the HLχQM, one encounters divergent loop integrals which will be quadratic, linear, and logarithmic divergent [12]. In order to obtain (15), (the regularized version of) these integrals has to be related to G_H, the pion coupling f_π, the constituent quark light mass m, and the gluon condensate. Similarly, in the light sector [13] the quadratic and logarithmic divergent integrals are related to the quark condensate and f_π, respectively.

To calculate the factorized contributions in (6) and (7) corresponding to Fig. 2 within our framework, we need the bosonized currents in (18) and (19), and in addition the bosonized currents involving emission of a photon from the B- or the D-meson. For photon emission from the B-meson we have (for $v = v_b$)

$$(\overline{q_L}\gamma^\alpha Q_{bv}^{(+)})_\gamma \to -\frac{G_H e}{32\pi} F_{\mu\nu} \mathrm{Tr}\left[\xi^\dagger \gamma^\alpha L H_{qb}^{(+)} \tilde{Q}_\xi \times \left(\sigma^{\mu\nu} - \frac{2\pi f_\pi^2}{m^2 N_c}\{\sigma^{\mu\nu}, \gamma\cdot v\}\right)\right], \qquad (23)$$

where F and $\tilde{Q}_\xi$ are given as in Eq. (17). For emission from the D-meson there is a similar expression. This structure is quite general, although the coefficients in front of the two terms are model dependent. A chiral Lagrangian piece corresponding to the diagram in Fig. 2(a) will be the product of this expression and one of the currents in Eq. (19).

Bosonizing currents with one gluon emission from a colored current occurring in the operators of (11) for instance to be used in the left part in Fig. 3(b)], one obtains

$$(\overline{q_L}t^a\gamma^\alpha Q_{bv}^{(+)})_G \to -\frac{G_H g_s}{64\pi} G_{\mu\nu}^a \mathrm{Tr}\left[\xi^\dagger \gamma^\alpha L H_{bq}^{(+)} \times \left(\sigma^{\mu\nu} - \frac{2\pi f_\pi^2}{m^2 N_c}\{\sigma^{\mu\nu}, \gamma\cdot v\}\right)\right], \qquad (24)$$

where $G_{\mu\nu}^a$ is the octet gluon tensor, and $H_{bq}^{(+)}$ represents the heavy $\bar{B}_q$-meson fields. Similarly the (heavy) D- and $\bar{D}$-mesons are represented by $H_c^{(+)}$ and $H_{\bar{c}}^{(-)}$ corresponding to a heavy quark field $Q_{v_c}^{(+)}$ and heavy antiquark field $Q_{\bar{v}}^{(-)}$, respectively, where v_c and $\bar{v} = v_{\bar{c}}$ are the velocities of the c and $\bar{c}$ quarks, respectively. The symbol $\{,\}$ denotes the anticommutator.

For one gluon and one photon emission from the $\bar{B}_q$-meson appearing in left part in Fig. 3(a)], one obtains an expression of the form

$$(\overline{q_L}t^a\gamma^\alpha Q_{v_b}^{(+)})_{G\gamma} \to G_H g_s e F_{\mu\nu} G_{\sigma\rho}^a \mathrm{Tr}[\xi^\dagger \gamma^\alpha L H_{bq}^{(+)} \times \tilde{Q}_\xi R^{\mu\nu\sigma\rho}], \qquad (25)$$

where the tensor R contains products of Dirac matrices (originating from vertices and propagators with momentum integrated out). Multiplying the currents for each vertex, for instance those in Eqs. (24) and (25), and using the prescription,

$$g_s^2 G_{\mu\nu}^a G_{\alpha\beta}^a \to 4\pi^2 \left\langle \frac{\alpha_s}{\pi} G^2 \right\rangle \frac{1}{12}(g_{\mu\alpha}g_{\nu\beta} - g_{\mu\beta}g_{\nu\alpha}), \qquad (26)$$

we obtain the bosonized version for the operator $\tilde{Q}_1$ in Eqs. (10) and (11) (with one extra photon field) as the product of two traces. The expression may be simplified by using the Dirac algebra, but we do not enter these details here.

As a more simple example, multiplying (24) for a decaying b-quark with the corresponding expression for creation of a c (or $\bar{c}$)-quark, and using (26), we obtain $1/N_c$ suppressed HLχPT terms for the virtual $B \to D$ transitions. Using some Dirac algebra inside the traces, we retain some color suppressed chiral Lagrangian term proportional to the dominant term being proportional to the product of the current in Eq. (18) and one of the currents in Eq. (19). Thus, such terms contribute to the weak transition coefficient β_W' of Ref. [5]. Also, some new $1/N_c$ terms occur, for instance

$$\mathrm{Tr}\{\overline{H_c^{(+)}}\xi R\sigma^{\mu\nu}\} \times \mathrm{Tr}\{\sigma_{\mu\nu} L \xi^\dagger H_b^{(+)}\}, \qquad (27)$$

$$\mathrm{Tr}\{\overline{H_c^{(+)}}\xi R\sigma^{\mu\nu}\} \times \mathrm{Tr}\{\sigma_{\mu\nu}\gamma\cdot t L \xi^\dagger H_b^{(+)}\}, \qquad (28)$$

$$\mathrm{Tr}\{\overline{H_c^{(+)}}\xi R\gamma\cdot t\sigma^{\mu\nu}\} \times \mathrm{Tr}\{\sigma_{\mu\nu} L \xi^\dagger H_b^{(+)}\}, \qquad (29)$$

where $t \equiv v_c - v_b$. Such terms have also to be considered when calculating pole diagrams. Note that the two last terms vanish in the low velocity limit $\omega \equiv v_b \cdot v_c \to 1$.

IV. AMPLITUDES FOR $B \to \gamma D^*$

Considering simple quark diagrams only, we observe that in terms of the Wolfenstein parameter λ, the amplitudes for $B^- \to \gamma D_d^{*-}$ and $\overline{B_d^0} \to \gamma \overline{D^{*0}}$ are $\mathcal{O}(\lambda^4)$ and small. In contrast, the amplitude for $\overline{B_d^0} \to \gamma D^{*0}$ is $\mathcal{O}(\lambda^2)$, and is KM nonsuppressed compared to other $B \to \gamma D^*$ modes. For $q = s$, all the amplitudes are $\mathcal{O}(\lambda^3)$. However, strong interactions in terms of meson exchange might make this simple picture more complicated.

At mesonic level there are *pole diagrams* obtained from the bosonized currents in (18) and (19) combined with photon emission [16,17]. Here the $B^- \to D_{s,d}^{*-}$ factorized transitions are proportional to the favorable coefficient $a_f = (a_{2,4} + a_{1,3}/N_c) \sim 1$, while the nonfactorizable contributions proportional to $2a_{1,3}$ due to colored operators are relatively small. For the decays $B^- \to \gamma D_{d,s}^-$, we obtain from the pole diagrams in Fig. 4,

$$A_{\text{pole}}^{(+)} = -(C_4 + C_3/N_c)\alpha_H^2 \frac{e\beta M_B}{M_B^2 - M_D^2}(Q_q^B - \omega Q_{q'}^D), \tag{30}$$

where Q_q^B and $Q_{q'}^D$ are the charges of the light quarks within the B and D mesons, respectively. For the processes $B^- \to \gamma D_{d,s}^-$, we have $Q_q^B = 2/3$ and $Q_{q'}^D = -1/3$. To obtain a parity violating pole term, the intermediate heavy meson(s) must have positive parity [5]. We find that within our framework, $B^- \to \gamma D_d^-$ has a partial branching ratio $\simeq 2 \times 10^{-7}$ from the diagram in Fig. 4.

For the decays $\overline{B_{d,s}^0} \to \gamma \overline{D^{*0}}$, there is a delicate balance between different amplitudes, and it is hard to conclude anything within our framework. For the decays $B^- \to \gamma D_{d,s}^-$, the factorized contributions dominate, and from the diagrams in Fig. 2 alone we obtain

$$A_F^{(\pm)} = -(C_4 + C_3/N_c)\frac{eN_c G_H \alpha_H}{16\pi} Y^{(\pm)}, \tag{31}$$

where

$$Y^{(+)} = Q_q^B + Q_{q'}^D(1 + k - k\omega y), \tag{32}$$

which within our framework is roughly 4 times the pole contribution. For the parity violating case, one has

$$Y^{(-)} = -Q_q^B(1 + 2k) + Q_{q'}^D(1 + k\omega y). \tag{33}$$

We find that within our framework, $B^- \to \gamma D_d^-$ has a partial branching ratio $\simeq 5 \times 10^{-7}$ from the "direct" factorizable diagram in Fig. 2, i.e. slightly bigger than our pole contribution.

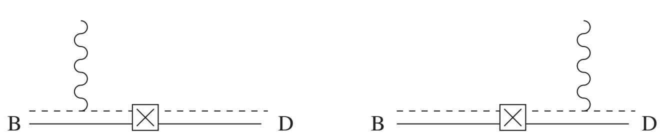

FIG. 4. Pole contributions for $B \to \gamma D$. The combined full and dashed lines are the heavy mesons, and the wavy lines represent photons.

As mentioned above, there are some *meson exchange* decay mechanism contributions. In the heavy quark limit these are identical to chiral loop contributions. These are shown in Fig. 5. For the process $B^- \to \gamma D_{s,d}^{*-}$ there is an intermediate $\overline{B_d^0} \to \overline{D^{*0}}$ transition accompanied with emission and reabsorption of π^-, or an intermediate $\overline{B_s^0} \to \overline{D^{*0}}$ transition accompanied with emission and reabsorption of K^-. For the process $\overline{B_d^0} \to \gamma \overline{D^{*0}}$ there is an intermediate $B^- \to D_d^{*-}$ transition accompanied with emission and reabsorption of π^+. For $\overline{B_s^0} \to \gamma \overline{D^{*0}}$ there is an intermediate $B^- \to D_s^{*-}$ transition accompanied with emission and reabsorption of K^+. Because the transitions $B^- \to D_{d,s}^{*-}$ are nonsuppressed in the factorized limit, the decays $\overline{B_{d,s}^0} \to \gamma \overline{D^{*0}}$ are semisuppressed, having meson exchange amplitudes reducing to chiral loops in the HQEFT limits $1/m_b \to 0$, and $1/m_c \to 0$. Taking both these limits, the meson exchange amplitudes are, in the leading logarithmic approximation, proportional to

$$\chi(M) = \left(\frac{g_A m_M}{4\pi f_\pi}\right)^2 \ln\left(\frac{\Lambda_\chi^2}{m_M^2}\right), \tag{34}$$

for exchange of light mesons $M = K, \pi$ respectively. Here g_A is the light meson axial coupling to heavy mesons and $\Lambda_\chi \simeq 1$ GeV. Numerically, $\chi(K) \simeq 0.09$ and $\chi(\pi) \simeq 0.02$. We have not estimated these meson exchange diagrams, neither in the low velocity limit as chiral loops nor in other ways (for instance, some form factor damping is expected). We just state that (34) should give the order of magnitude for meson exchanges. For the processes $\overline{B_{d,s}^0} \to \gamma D^{*0}$ there are only Zweig-forbidden and $SU(3)_F$ violating neutral meson exchange which give small contributions.

For the processes $\overline{B_{d,s}^0} \to \gamma D^{*0}$ the nonfactorizable amplitudes $A^{(\pm)}$ corresponding to the diagrams in Fig. 3 are of the form

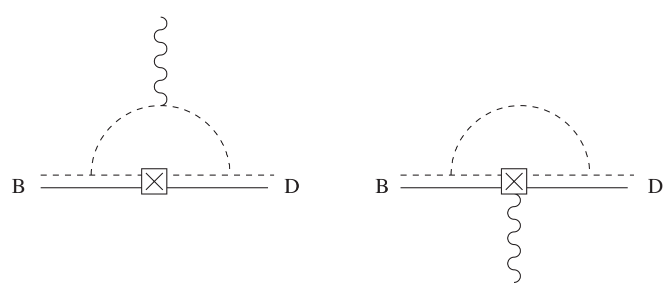

FIG. 5. Meson exchange diagrams which would be chiral loops in the low velocity limit. The combined full and dashed lines are the heavy mesons, and the single dashed lines represent light pseudoscalar mesons.

$$A_G^{(\pm)} = -\frac{eC_2}{2^8\pi} G_H^2 \left\langle \frac{\alpha_s}{\pi} G^2 \right\rangle (Q_q^B Z_B^{(\pm)} + Q_{q'}^D Z_D^{(\pm)}). \qquad (35)$$

For the processes $\overline{B^0_{d,s}} \to \gamma D^{*0}$ we have $Q_q^B = -1/3$ and $Q_{q'}^D = 2/3$. The quantities $Z^{(\pm)}$ are of order one and given by [22]

$$Z_B^{(+)} = \left(\frac{89\pi}{288} - \frac{13}{18}\right) k\omega y + \left(\frac{7\pi}{144} + \frac{5}{18}\right) k\omega^2 + \left(\frac{11}{18} - \frac{13\pi}{96}\right) k + \left(\frac{2}{3} - \frac{\pi}{18}\right), \qquad (36)$$

$$Z_B^{(-)} = -\frac{5\pi}{9} k\omega y + \frac{(\pi+2)}{9} k\omega^2 + \left(\frac{\pi}{9} - \frac{4}{3}\right) k - \frac{(\pi+8)}{9}, \qquad (37)$$

$$Z_D^{(+)} = -\left(\frac{11\pi}{288} + \frac{17}{36}\right) k\omega y + \left(\frac{1}{36} - \frac{53\pi}{288}\right) k - \left(\frac{\pi}{64} + \frac{7}{72}\right) \omega y + \left(\frac{41\pi}{576} + \frac{23}{72}\right), \qquad (38)$$

$$Z_D^{(-)} = -\left(\frac{\pi}{3} + \frac{4}{9}\right) k + \left(\frac{2}{9} - \frac{\pi}{18}\right) \omega y + \frac{4}{3} k\omega y + \left(\frac{2}{9} - \frac{\pi}{18}\right), \qquad (39)$$

where the dimensionless parameters k, ω, and y are defined as

$$k = \frac{2\pi f_\pi^2}{N_c m^2}, \qquad \omega = v_b \cdot v_c = \frac{M_B^2 + M_D^2}{2M_B M_D}, \qquad y = \frac{M_B}{M_D}. \qquad (40)$$

For M_D we have used the mass of D^*. Using [12,17] $m = 230$ MeV, $\rho = 1.1$, and $\langle \frac{\alpha_s}{\pi} G^2 \rangle^{1/4} = 310$ MeV, we obtain

$$\mathrm{BR}(\overline{B^0_d} \to \gamma D^{*0}) \simeq 1 \times 10^{-5} \quad \text{and} \quad \mathrm{BR}(\overline{B^0_s} \to \gamma D^{*0}) \simeq 6 \times 10^{-7}. \qquad (41)$$

V. CONCLUSION

We have considered and classified six decay modes of the type $B \to \gamma D$ generated by three (main) mechanisms:

(a) Factorized contributions of pole and nonpole type. These are proportional to the favorable Wilson coefficient combination $a_f \sim 1$ for decays of charged B-mesons. However, for decays of neutral B-mesons, they are proportional to the nonfavorable coefficient combination a_{nf} of order 10^{-2}, and are *unimportant*.

(b) Nonfactorizable contributions due to *meson exchanges*, that is, some intermediate $B \to D$ transition accompanied with an emission and reabsorption of a pseudoscalar boson (π or K). Such contributions are suppressed (Zweig-forbidden) for $\overline{B^0_{d,s}} \to \gamma D^{*0}$.

(c) Nonfactorizable contributions due to the "colored quark operators" in Eqs. (10) and (11), obtained from Fierz transformations of the standard four quark operators. Within the HLχQM, these are calculated as in terms of emission of soft gluons, forming a (model dependent) gluon condensate. These contributions are proportional to the small Wilson coefficients $a_{1,3}$, and play a minor role for decays of charged B-mesons. For decays of neutral B-mesons, they are proportional to the favored coefficients $a_{2,4} \sim 1$, and are *important*, and for $\overline{B^0_{d,s}} \to \gamma D^{*0}$ they dominate.

In the heavy quark effective field theory (HQEFT) limits, the mechanisms (b) and (c) are formally $1/N_c$ suppressed. However, they play an important role for the neutral decay modes. Contributions to the various $B \to \gamma D^*$ modes are classified in Table I.

The present analysis is performed within HQEFT and heavy-light chiral perturbation theory (HLχPT), both for the b- and the c-quark. Formally, the modes $B \to \gamma D^*$ are "heavy to heavy" transitions, but for precise estimates our framework is not ideal [4] because the energy gap between the b- and the c-quark are significantly bigger than 1 GeV, which is the scale of HLχPT and the heavy-light chiral quark model (HLχQM). Therefore large $1/m_Q$ corrections (especially $1/m_c$ corrections) must be expected. Phrased in another way, damping form factors are expected to be present, and our estimates might be overestimates. A very recent study [5] shows that positive parity resonances in pole diagrams might even give slightly bigger amplitudes than obtained in the present paper. The difference between the paper [5] and the present paper is that we do not consider positive parity intermediate resonances. But they will anyway not contribute to the nonfactorizable amplitudes which are the main task of this paper. More important is that in [5] the Wilson coefficients at the renormalization scale $\mu \sim M_W$ are used, while we have used the Wilson coefficients at the matching scale $\mu \sim \Lambda_\chi \simeq 1$ GeV. This has dramatic consequences for the color suppressed amplitudes. We have an effective Wilson coefficient a_{nf} of order 10^{-2} for the factorized amplitudes, while Ref. [5] uses an effective Wilson coefficient $a_{nf} = 1/N_c = 1/3$. With this big value of a_{nf}, the authors of [5] obtain pole contributions for $\overline{B^0_{d,s}} \to \gamma D^{*0}$ of the same

TABLE I. A classification of contributions to processes of the type $B \to \gamma D^*$.

Process	Factorized	Soft gluon	Meson exchange
$B^- \to \gamma D_d^{*-}$	$a_f \lambda^4$	$a_3 \lambda^4$	$a_{nf} \lambda^4 \chi(\pi^-)$
$B^- \to \gamma D_s^{*-}$	$a_f \lambda^3$	$a_3 \lambda^3$	$a_{nf} \lambda^3 \chi(K^-)$
$\overline{B^0_d} \to \gamma \overline{D^{*0}}$	$a_{nf} \lambda^4$	$a_4 \lambda^4$	$a_f \lambda^4 \chi(\pi^+)$
$\overline{B^0_s} \to \gamma \overline{D^{*0}}$	$a_{nf} \lambda^3$	$a_4 \lambda^3$	$a_f \lambda^3 \chi(K^+)$
$\overline{B^0_d} \to \gamma D^{*0}$	$a_{nf} \lambda^2$	$a_2 \lambda^2$	Zweig-forbidden
$\overline{B^0_s} \to \gamma D^{*0}$	$a_{nf} \lambda^3$	$a_2 \lambda^3$	Zweig-forbidden

order of magnitude as we obtain from the colored quark operators.

Alternative estimates, based on other frameworks, for instance considering the charm quark as "light," might be performed. Still, we expect that we have obtained *amplitudes* of the right order of magnitude for $\overline{B^0_{d,s}} \to \gamma D^{*0}$. Namely, the matrix element of the colored operators in (10) and (11) are in general nonzero and give an important contribution to $\overline{B^0_{d,s}} \to \gamma D^{*0}$.

It should be noted that our framework will be better suited to describe the processes $B \to D\gamma\gamma$ [5], where ω might be as low as $\simeq 1.2$ if the two photons come out back to back, in contrast to $\omega \simeq 1.6$ in our case.

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