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$B^+ \rightarrow D_s^{*+}\gamma$ and $B^+ \rightarrow D^{*+}\gamma$ as Probes of V_{ub}

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Abstract

The decays $B^+ \rightarrow D_s^{*+}\gamma$ and $B^+ \rightarrow D^{*+}\gamma$ can be used for an extraction of $|V_{ub}|$. When the b and c quarks are nearly degenerate the rate for these modes can be determined in terms of other observed rates, namely $B\bar{B}$ mixing and $D^* \rightarrow D\gamma$ decay. To this end we introduce a novel application of heavy quark and flavor symmetries. Although somewhat unrealistic, this limit provides us with a first estimate of these rates.

12.15.Hh, 12.39.Hg, 13.40.Hq, 13.25.Hw

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I. INTRODUCTION

The extraction of the Cabibbo-Kobayashi-Maskawa (CKM) mixing matrix element V_{ub} , theoretically and experimentally, stands out as one of the prominent challenges of particle physics. Both its magnitude (through decay rates) and its phase relative to other CKM elements (through CP violation) are of considerable interest. Nevertheless, its determination remains elusive, ultimately because the CKM matrix describes the mixing of quarks, whereas of course only hadrons are observed. In the case of V_{cb} , much of this impediment is overcome by the application of the heavy quark effective theory (HQET) [1,2], which in particular relates the strong interaction matrix elements of systems with b and c quarks, eliminating much of the strong interaction uncertainty. However, the u quark is by no means heavy, and it is notoriously difficult to separate its CKM and strong interaction couplings.

This is true even for the experimentally clean semileptonic modes, $\bar{B} \rightarrow \pi \ell \bar{\nu}$ and $\bar{B} \rightarrow \rho \ell \bar{\nu}$, although dispersive bounds on the shapes of the form factors help to restrict such strong interaction uncertainties [3]. Inclusive semileptonic rates for $\bar{B} \rightarrow X_u \ell \bar{\nu}$ are theoretically under better control but are plagued, since $|V_{cb}| \gg |V_{ub}|$, by the preponderance of $\bar{B} \rightarrow X_c \ell \bar{\nu}$ everywhere in kinematic space except in the endpoint region of maximal lepton energy, where a c quark cannot kinematically be produced; however, in this small region it has proved necessary to include hadronic model dependence [4]. A technique involving invariant hadronic mass spectra for invariant mass below m_D is under better theoretical control but requires neutrino reconstruction [5].

Hadronic modes, containing everywhere strong interaction uncertainties, are even more problematic. Even if these difficulties could be tamed, one would still encounter the mixing of different weak topologies. For example, $\bar{B}^0 \rightarrow \pi^+ K^-$ may proceed through either $b \rightarrow us\bar{u}$ ($\propto V_{ub}V_{us}^*$) or a penguin $b \rightarrow sg \rightarrow s\bar{u}u$ ($\propto V_{tb}V_{ts}^*$); the presence of spectator quarks complicates the analysis.

One may probe V_{ub} using the decays $B^+ \rightarrow D_s^{*+} \gamma$ ($\bar{b}u \rightarrow c\bar{s}\gamma$) and $B^+ \rightarrow D^{*+} \gamma$ ($\bar{b}u \rightarrow c\bar{d}\gamma$) [6], collectively $B^+ \rightarrow D_{(s)}^{*+} \gamma$. Even though the second decay is Cabibbo-suppressed compared to the first, the lower reconstruction efficiency associated with D_s^{*+} compared to D^{*+} makes both processes worth studying. The hard, monochromatic photon in these decays, 2.22 and 2.26 GeV, respectively, provides a distinctive experimental signature.

As shown in Fig. 1, the invariant amplitude for these decays at $O(G_F^1)$ consists of only one weak topology, because four flavor quantum numbers change ($\Delta B = -1$, $\Delta C = +1$, $\Delta(S, D) = +1$, $\Delta U = -1$), thus fixing the quark couplings to the W boson.¹ The photon may couple to any of the quark lines, although one expects the dominant contributions from radiation emitted by the light u or \bar{s} . The diagram with the photon emitted from the W is suppressed, of $O(G_F^2)$. On the other hand, since gluons may pair-produce quarks, these decays may be sensitive to multiparticle intermediate states, for example $B^+ \rightarrow D^0 K^+ \rightarrow D_s^{*+} \gamma$ [7], but the weak topology remains unchanged.

In this paper we study a theoretical limit in which the decays $B^+ \rightarrow D_{(s)}^{*+} \gamma$ are calculable from first principles in terms of other measured quantities. We assume both the b and c

quarks are heavy and nearly degenerate, $\Lambda_{\text{QCD}} \ll m_b$, $\Lambda_{\text{QCD}} \ll m_c$ and $\delta m \equiv m_b - m_c \lesssim \Lambda_{\text{QCD}}$. We do not expect this limit to be a good approximation to reality. Our calculation is a starting point for further investigations on controlled approximations of this rate.

¹At $O(G_F^2)$, the nontrivial diagrams include the di-penguin $\bar{b} \rightarrow \bar{s}g$, $ug \rightarrow c$ and a box diagram with the internal quarks crossed.

FIGURES

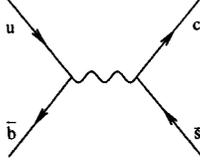


FIG. 1. The unique weak diagram topology at $O(G_P^1)$ for $B^+ \rightarrow D_s^{*+}\gamma$. The diagram for $B^+ \rightarrow D^{*+}\gamma$ simply changes \bar{s} to \bar{d} .

II. FORMALISM

The dominant contributions to the amplitude for $B^+ \rightarrow D_{(s)}^{*+}\gamma$ are the long distance processes $B^+ \rightarrow D_{(s)}^+ \rightarrow D_{(s)}^{*+}\gamma$ and $B^+ \rightarrow B^{*+}\gamma \rightarrow D_{(s)}^{*+}\gamma$. The premise of this calculation is that the matrix elements of $B^+ \rightarrow D_{(s)}^+$, $D_{(s)}^+ \rightarrow D_{(s)}^{*+}\gamma$, and $B^+ \rightarrow B^{*+}\gamma$ can be related by heavy quark and $SU(3)$ flavor symmetries to the rates for $B\bar{B}$ mixing and $D_{(s)}^* \rightarrow D_{(s)}\gamma$. Alternately, we expect that lattice calculations similar to those that determine the $B\bar{B}$ mixing matrix element [8] can be employed to compute that of $B^+ \rightarrow D_{(s)}^+$ directly.

For the initial calculation presented here, we work in the generally-low (GL) velocity limit

$$\delta m \equiv m_b - m_c \lesssim \Lambda_{\text{QCD}} \ll m_b, \quad (2.1)$$

which means that the four-velocity of each heavy particle in the process is assumed to remain constant to lowest order. This limit is more restrictive than the slow-velocity (SV) limit considered by Shifman and Voloshin [2], $\Lambda_{\text{QCD}} \ll \delta m \ll m_b$. To indicate the quality of this assumption, we point out that the value of $\gamma_{D_{(s)}^*} = v_B \cdot v_{D_{(s)}^*}$ is 1.45 for the strange case and 1.50 for the nonstrange case. On the other hand, letting $m_B = m_b + \Lambda - 3\lambda^2/4m_b$, $m_{D_{(s)}^*} = m_c + \Lambda + \lambda^2/4m_c$, one finds that $\gamma_{D_{(s)}^*}$ is parametrically small,

$$\begin{aligned} \gamma_{D_{(s)}^*} &= \frac{m_B^2 + m_{D_{(s)}^*}^2}{2m_B m_{D_{(s)}^*}} \\ &= 1 + \frac{1}{2m_b^2} \delta m^2 + O\left(\frac{1}{m_b^3}\right). \end{aligned} \quad (2.2)$$

Heavy quark symmetries may be used to relate the matrix element for $B \rightarrow D_{(s)}$ to the one for $B \rightarrow \bar{B}$ only if the velocity of the $D_{(s)}$ in the B rest frame is parametrically small, say $O(\Lambda/m_b)$ or $O(\delta m/m_b)$. It is therefore sufficient to assume the weaker SV limit for this part of our argument. However, in order to relate the photon emission processes to the real decays $B^* \rightarrow B\gamma$ and $D^* \rightarrow D\gamma$, it is necessary to impose the GL condition to ensure softness of the photon, as we now argue.

The matrix element of the electromagnetic current j_μ between vector (D^* or B^*) and pseudoscalar (D or B) states is characterized by a single Lorentz invariant form factor

$$\langle \bar{p}', \bar{\epsilon} | j_\mu(0) | \bar{p} \rangle = -ig(q^2) \epsilon_{\mu\nu\lambda\sigma} \epsilon^{*\nu} p'^\lambda p^\sigma, \quad (2.3)$$

where $q = (p - p')$. The form factor at $q^2 = 0$ governs the on-shell decay of the vector to the pseudoscalar, which is characterized by a transition magnetic moment, $g(0) = e\mu$ [see below, Eq. (3.1)]. However, implicit in this definition is that the mesons are on-shell states; if the transition occurs with one of the states being virtual, as is the case here, then the q^2 argument of the form factor must be computed accordingly. The relevant scale over which a hadronic form factor changes is Λ_{QCD} . For example, in the non-relativistic potential quark model of mesons, the scale that dictates the behavior of the form factor g is the constituent quark mass $\sim \Lambda_{\text{QCD}}$. In $B^+ \rightarrow D_{(s)}^{*+}\gamma$ the form factor appears with argument $q^2 \simeq -\delta m^2$, and in the GL limit one may approximate $g(-\delta m^2) \simeq g(0) = e\mu$.

The formalism of HQET combined with chiral symmetry, relevant to the physics of mesons containing heavy quarks, was developed in Refs. [9]. It was extended to radiative processes, such as $D^* \rightarrow D\gamma$, in Ref. [10], while the Lagrangian for heavy quark-heavy antiquark mixing appeared in [11]. We use ingredients from all of these works, but for brevity include only notation relevant to the present process. While convenient, the formalism is not really a necessary framework for our calculation. In fact, light mesons enter diagrams only at higher order in the chiral expansion. By stating our calculation in the language of this effective Lagrangian we are setting the stage for further investigations of, for example, the effects of the finite strange quark mass.

Since chiral symmetry is one of the ingredients of the Lagrangian, we begin by including the octet of pseudo-Nambu-Goldstone bosons in the usual nonlinear form ξ :

$$\xi = \exp(i\Pi/f), \quad (2.4)$$

where

$$\Pi = \begin{pmatrix} \frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta & \pi^+ & K^+ \\ \pi^- & -\frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta & K^0 \\ K^- & \bar{K}^0 & -\sqrt{\frac{2}{3}}\eta \end{pmatrix}, \quad (2.5)$$

with $f \simeq 131$ MeV. Under the chiral $SU(3)_L \times SU(3)_R$ transformation (L, R),

$$\xi \mapsto L\xi U^\dagger = U\xi R^\dagger, \quad (2.6)$$

with U implicitly defined by the equality of these two forms.

The ground state vector P_{Aa}^* and pseudoscalar P_{Aa} fields that destroy mesons of heavy quark flavor A and light antiquark flavor a are incorporated into the 4×4 bispinor

$$\begin{aligned} H_{Aa} &= \frac{(1 + \not{v})}{2} (P_{Aa}^{*\mu} \gamma_\mu - P_{Aa} \gamma_5), \\ \bar{H}^{Aa} &= \gamma^0 H_{Aa}^\dagger \gamma^0. \end{aligned} \quad (2.7)$$

Under $SU(4)$ heavy quark spin-flavor symmetry transformations S and $SU(3)_L \times SU(3)_R$ transformations U , H_a transforms as

$$\begin{aligned} H_{Aa} &\mapsto S_A^B H_{Bb} U^{\dagger b}{}_a, \\ \bar{H}^{Aa} &\mapsto U^a{}_b \bar{H}^{Bb} S^{\dagger B}{}_A. \end{aligned} \quad (2.8)$$

Defining the chiral covariant derivative and axial current by

$$\begin{aligned} D_{ab}^\mu &= \delta_{ab} \partial^\mu - V_{ab}^\mu \\ &= \delta_{ab} \partial^\mu - \frac{1}{2} (\xi^\dagger \partial^\mu \xi + \xi \partial^\mu \xi^\dagger)_{ab} = \delta_{ab} \partial^\mu + O(\Pi^2), \\ A_{ab}^\mu &= \frac{i}{2} (\xi^\dagger \partial^\mu \xi - \xi \partial^\mu \xi^\dagger)_{ab} = -\frac{1}{f} \partial^\mu \Pi_{ab} + O(\Pi^3), \end{aligned} \quad (2.9)$$

one finds that $D\xi$, $D\xi^\dagger$, DH , and $D\bar{H}$ have precisely the same transformation properties under chiral symmetry as their underlying fields, and

$$(A^\mu)^a{}_b \mapsto U^a{}_c (A^\mu)^c{}_d U^{\dagger d}{}_b. \quad (2.10)$$

Thus, one obtains the heavy meson Lagrangian invariant under chiral and heavy quark symmetry with the minimum number of derivatives,

$$\mathcal{L} = -\text{Tr} \bar{H}^{Aa} i v \cdot D^b{}_a H_{Ab} + g \text{Tr} \bar{H}^{Aa} H_{Ab} A^b{}_a \gamma_5. \quad (2.11)$$

It should be noted that the kinetic term is canonically normalized when the field H has 2, rather than the usual $2M$, particles per unit volume.

Since lowest-order HQET integrates out heavy antiparticle degrees of freedom, it is necessary to include such fields explicitly when they can appear in the asymptotic states. One defines [11]

$$\begin{aligned} H^{\bar{A}a} &= \left(P_\mu^* \bar{A}^a \gamma^\mu - P^{\bar{A}a} \gamma_5 \right) \frac{(1 - \not{v})}{2}, \\ \bar{H}_{\bar{A}a} &= \gamma^0 H^{\dagger \bar{A}a} \gamma^0. \end{aligned} \quad (2.12)$$

For example, while the field P_{Aa}^* destroys a vector meson of flavor content $A\bar{a}$, the field $P^{\bar{A}a}$ destroys one of flavor content $\bar{A}a$; the two are related by a chosen charge conjugation convention [11]. One demands that H_{Aa} and $H^{\bar{A}a}$, or \bar{H}^{Aa} and $\bar{H}_{\bar{A}a}$, have the same transformation properties under heavy quark symmetry and chiral transformations. Then the construction of the Lagrangian for mesons with heavy antiflavors is straightforward.

As for direct mixing between the two sectors, the standard model operator that destroys quarks of flavors \bar{A} , a and creates quarks of flavors B , \bar{b} , given charge constraints $-Q_A + Q_a = Q_B - Q_b$, is

$$\mathcal{O}^{ab, \bar{A}B} = \bar{A} \gamma^\mu (1 - \gamma_5) a \cdot \bar{B} \gamma_\mu (1 - \gamma_5) b. \quad (2.13)$$

For $B\bar{B}$ mixing, one sees that A and B are b quarks, and a and b are d quarks; in the present case, A , B , a , and b are b , c , u , and $(d$ or $s)$ quarks, respectively. In the limit that HQET and SU(3) are good symmetries, both processes have the same strong matrix element. In Ref. [11] it was shown that (2.13) matches in the symmetry limit of the effective theory onto

$$\mathcal{O}^{ab, \bar{A}B} = 4\beta \left[(\xi P^{*+i\mu})^{\bar{A}a} (\xi P^*)^{Bb} + (\xi P^{\dagger})^{\bar{A}a} (\xi P)^{Bb} \right] + \dots \quad (2.14)$$

Calculation of the process also requires understanding the coupling of heavy mesons to photons. Minimal substitution for the kinetic term in Eq. (2.11) provides no coupling to real transverse photons at lowest order, while minimal substitution applied to the second term in (2.11) requires including at least one pion for a nonvanishing term. The unique lowest-order electromagnetic term [10] is the transition magnetic moment operator

$$\delta\mathcal{L} = \frac{1}{4} e \mu_{Aa} \text{Tr} \bar{H}^{Aa} H_{Aa} \sigma^{\mu\nu} \mathcal{F}_{\mu\nu}^{\text{EM}} + \text{charge conjugate}. \quad (2.15)$$

III. CALCULATION

The photon coupling in Eq. (2.15) represents the lowest-order heavy-to-heavy electromagnetic transition. It does not couple pseudoscalars to each other, but provides for the decay of vector mesons:

$$\Gamma(P_{Aa}^* \rightarrow P_{Aa} \gamma) = \frac{1}{3} \alpha |\mu_{Aa}|^2 E_\gamma^3. \quad (3.1)$$

To lowest order in HQET and SU(3), one may write² $\mu_{Aa} = Q_a \mu$. For $P_{Aa}^* = D^{*+}$, the current experimental bound [12] of $\Gamma(D^{*+}) < 0.131$ MeV and electromagnetic branching fraction $< 3.2\%$, lead to $|\mu| < 2.5$ GeV⁻¹.

The calculation of the rate is then straightforward in the combined HQET and chiral symmetry limit, and consists of the two diagrams depicted in Fig. 2. We use the GL limit to compute the amplitude, but retain full mass dependence in the phase space calculation. The result is

$$\Gamma(B^+ \rightarrow D_s^{*+} \gamma) = \frac{G_F^2}{64} |V_{ub}^* V_{cs}|^2 (4\beta)^2 \alpha |\mu|^2 \frac{m_{D_s^*}}{m_B^4} (m_B - m_{D_s^*})(m_B + m_{D_s^*})^3, \quad (3.2)$$

with $V_{cs} \rightarrow V_{cd}$ and $m_{D_s^*} \rightarrow m_{D^*}$ for $B^+ \rightarrow D^{*+} \gamma$.



FIG. 2. The two leading diagrams for $B^+ \rightarrow D_s^{*+} \gamma$ that occur in the heavy quark and chiral limits. The double line indicates a heavy meson, one containing a b (c) quark on the left (right) of the flavor-changing vertex, indicated by a box.

²In Eq. (14) of Ref. [10], μ is labeled β .

The effective Lagrangian mixing parameter β in (2.14) is related to familiar quantities of $B\bar{B}$ mixing via

$$\frac{1}{m_B} \langle \bar{B}^0 | \mathcal{O}^{du, \bar{B}B} | B^0 \rangle = \frac{8}{3} f_B^2 m_B B_B = 4\beta. \quad (3.3)$$

A full treatment requires including renormalization point dependence in B_B , but we neglect this here for simplicity. Now there are two obvious directions of analysis. One is to eliminate β using its calculated value from lattice simulations. Using [8] $f_B = 170 \pm 35$ MeV and $B_B = 0.98 \pm 0.06$ (renormalization point 2 GeV), and $m_B = 5.279$ GeV, $\tau_B = 1.6 \cdot 10^{-12}$ s, one finds $4\beta = 2.35 \pm 0.50$ GeV³, and

$$\text{Br}(B^+ \rightarrow D_s^{*+} \gamma) \simeq 4 \cdot 10^{-5} \cdot \left(\frac{B_B}{0.98} \right)^2 \cdot \left(\frac{\Gamma(D^{*+})}{0.131 \text{ MeV}} \right) \cdot \left(\frac{\text{Br}(D^{*+} \rightarrow D^+ \gamma)}{3.2\%} \right) \cdot \left| \frac{V_{ub}^* V_{cs}}{3 \cdot 10^{-3}} \right|^2, \quad (3.4)$$

with $V_{cs} \rightarrow V_{cd}$ for $B^+ \rightarrow D^{*+} \gamma$, in which case the coefficient (taking $|V_{cd}| = 0.22$) is $2 \cdot 10^{-6}$. The other direction for analysis eliminates β as a systematic strong interaction uncertainty between this process and $B\bar{B}$ mixing. In particular, $\Delta m_B \propto G_F^2 |V_{ub}^* V_{td}^*|^2 \beta$, which means³ that eliminating β leads to values for the ratio $|V_{ub}^* V_{cs}|^2 / |V_{ub}^* V_{td}^*|^4$, which is proportional to $(\rho^2 + \eta^2) / [(1-\rho)^2 + \eta^2]^2$ in the usual Wolfenstein space. Curves of this family tend to intersect the ρ axis with vertical slope, which is useful since the current experimentally allowed region in ρ - η space is broad in the ρ direction.

IV. PROSPECTS

The inclusion of HQET-violating, chiral symmetry-violating, and short-distance corrections is straightforward, and this program should certainly be carried out. Indeed, many of the necessary corrections already appear in the literature, including corrections to μ_{Aa} [10], meson decay constants [11,13], $B\bar{B}$ mixing [11], and a number of other SU(3)-symmetry corrections [11,14]. However, a number of other diagrams remain to be computed, because the inclusion of the photon changes the allowed quantum numbers of intermediate states. For example, there are new chiral loop diagrams for $BD_{(s)}$ mixing compared to $B\bar{B}$ mixing where a single pion emerges directly from the mixing vertex.

More urgent and less tractable is understanding of the radiative form factor $g(q^2)$ at realistic momenta. Taking $g(q^2) \rightarrow g(0) = e\mu$ introduces large uncertainties into our estimate for the decay rate. Alternatively, one may use only the less extreme SV limit to relate the mixing amplitudes and express the rate in terms of the corresponding form factors. This is done by replacing $e\mu$ in Eq. (3.2) by $g_B(q^2) + g_D(q^2)$ at $q^2 = -\delta m^2$, which can be used as a starting point for a calculation based on estimates using hadronic models for the form factors.

³Note that $\Gamma(B^+ \rightarrow D_{(s)}^{*+} \gamma) \propto G_F^2 \beta^2$, while $\Delta m_B \propto G_F^2 \beta$, due to the different weak topologies.

We expect that the best determination of V_{ub} will still require careful analysis of a number different decay modes; however, the decays $B^+ \rightarrow D_{(s)}^{*+} \gamma$ provide an important additional handle on the problem.

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