Radiative heavy meson transitions

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ABSTRACT

We evaluate the radiative and hadronic decay rates of the $D^*$ mesons using the Heavy Quark Effective Theory and the Vector Meson Dominance hypothesis. We also estimate the width of the $B^*$ electromagnetic transitions and the radiative decays of positive parity $J^P = 0^+, 1^+$ charmed mesons.

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In this letter we wish to show that the heavy quark and chiral \( SU(3)_L \times SU(3)_R \)
symmetries of QCD, together with the Vector Meson Dominance (VMD) hypothesis, can be used to relate radiative and semileptonic charmed meson decays. We shall show that, using experimental data on the \( D \to \pi \ell \nu_\ell, D \to K \ell \nu_\ell \) and \( D \to K^* \ell \nu_\ell \) transitions as an input, it is possible to predict the radiative branching ratios \( D^* \to D\gamma \), the hadronic rate \( D^* \to D\pi \) and therefore the full \( D^* \) width. Moreover, the heavy quark flavour symmetry will provide us with a prediction of the \( B^* \) decay rate. We shall also employ information on the positive parity \( J^P = 0^+, 1^+ \) charmed mesons (masses and couplings) from the analysis of semileptonic \( D \) decays via axial currents in order to get predictions on the radiative decays of these states as well.

The calculation is based on a chiral lagrangian approach to the heavy and light meson interactions incorporating the chiral symmetry and the heavy quark spin flavour symmetry \([1, 2]\). At the lowest order in the light meson derivatives, the chiral lagrangian can be written as follows:

\[
\mathcal{L} = \mathcal{L}_0 + \mathcal{L}_2 + \ldots ,
\]

where \( \mathcal{L}_0 \) contains the light meson matrix \( \Sigma \) and the heavy meson fields \( H \):

\[
\mathcal{L}_0 = \frac{f^2_\pi}{8} < \partial^\mu \Sigma \partial_\mu \Sigma^\dagger > + i < H_b V^\mu D_{\mu ba} \bar{H}_a > + ig < H_b \gamma_\mu \gamma_5 A^\mu_{ba} \bar{H}_a > ,
\]

whereas \( \mathcal{L}_2 \) \([2]\) describes the interactions with the light \( J^P = 1^- \) mesons:

\[
\mathcal{L}_2 = - \frac{f^2_\pi}{2} a < (V_\mu - \rho_\mu)^2 > + \frac{1}{2g_V^2} < F_{\mu\nu}(\rho) F^{\mu\nu}(\rho) > \\
+ i \beta < H_b v^\mu (V_\mu - \rho_\mu)_{ba} \bar{H}_a > \\
+ \frac{\beta^2}{2f^2_\pi a} < \bar{H}_b H_a H_a H_b > + i \lambda < H_b \sigma^{\mu\nu} F_{\mu\nu}(\rho)_{ba} \bar{H}_a > .
\]

In eqs.\([2,3]\) \( < \ldots > \) means the trace, \( f_\pi = 132 \) MeV and:

\[
D_{\mu ba} = \delta_{ba} \partial_\mu + V_{\mu ba} = \delta_{ba} \partial_\mu + \frac{1}{2} \left( \xi^\dagger \partial_\mu \xi + \xi \partial_\mu \xi^\dagger \right)_{ba} ,
\]

\[
A_{\mu ba} = \frac{1}{2} \left( \xi^\dagger \partial_\mu \xi - \xi \partial_\mu \xi^\dagger \right)_{ba} ;
\]
the field $\xi$ is defined by:

$$\xi = \sqrt{\Sigma} = e^{iM/f_*}$$  \hfill (6)

where $M_{ba}$ is the usual $3 \times 3$ matrix describing the octet of pseudo Nambu-Goldstone bosons. The $0^-(P)$ and $1^-(P^*)$ $Q\bar{q}_a$ heavy mesons are described by the effective fields:

$$H_a = \frac{(1 + \not{v})}{2} [P^*_{a\mu} \gamma^\mu - P_a \gamma_5],$$  \hfill (7)

$$\bar{H}_a = \gamma_0 H_a^\dagger \gamma_0,$$  \hfill (8)

where $v$ is the heavy meson four-velocity, $a = 1, 2, 3$ (for $u, d$ and $s$ respectively), and $P^*_{a\mu}$ and $P_a$ are annihilation operators normalized as follows:

$$\langle 0|P_a|Q\bar{q}_a(0^-)\rangle = \sqrt{m_P};$$

$$\langle 0|P^*_{a}\bar{q}_a(1^-)\rangle = e^\mu \sqrt{m^*_P}$$

(with $m_P = m_{P^*}$ in the limit of heavy quark spin symmetry).

The coupling of the light vector meson resonances belonging to the $1^-$ low lying nonet (with $\phi = s\bar{s}$) has been introduced in $L_2$ (eq.3) using the hidden gauge symmetry approach $^3$; $\rho^\mu$ contains the light $1^-$ fields, normalized according to:

$$\rho^\mu = i \frac{g_V}{\sqrt{2}} \hat{\rho}^\mu$$  \hfill (9)

where $\hat{\rho}$ is the $3 \times 3$ matrix describing the $1^-$ light meson nonet; $F_{\mu\nu}(\rho) = \partial_\mu \rho_\nu - \partial_\nu \rho_\mu + [\rho_\mu, \rho_\nu]$ and $g_V$, $a$ and $g$ are coupling constants. The two KSRF relations $^4$ fix the value of $a$ and $g_V$:

$$a = 2 \hspace{2cm} g_V = 5.8.$$  \hfill (10)

The Lagrange density $^1$ is the most general expression displaying chiral invariance in the lowest order in the meson derivatives; this invariance can be proved observing that under $SU(3)_L \times SU(3)_R$ transformations:
\[ \xi \rightarrow L \xi U^\dagger = U \xi R^\dagger, \]
\[ \mathcal{A}^\mu \rightarrow U \mathcal{A}^\mu U^\dagger, \]
\[ D^\mu \bar{H} \rightarrow UD^\mu \bar{H}. \]

The Lagrangian (1) displays also the heavy quark flavour symmetry (since the mass of the heavy quark does not appear in \( \mathcal{L} \)) and the heavy quark spin symmetry, because, under the heavy quark spin group \( SU(2)_v \), the fields transform as follows:

\[ H_a \rightarrow SH_a; \quad \bar{H}_a \rightarrow \bar{H}_a S^\dagger, \]

with \( SS^\dagger = 1 \) and \([S, S] = 0\).

Let us now discuss the constants \( g \) and \( \lambda \) (\( \beta \) will not be used hereafter); \( g \) is responsible for the \( D^*D\pi \) coupling and is related to the \( D^* \) hadronic width by the tree level formula:

\[ \Gamma (D^* \rightarrow D^0 \pi^+) = \frac{g^2}{6\pi f^2_\pi} |\vec{p}_\pi|^3; \]

\((\Gamma (D^* \rightarrow D^+ \pi^0) \) is smaller by a factor of 2). As shown in [2], the coupling \( g \) can be obtained by the \( D \rightarrow \pi \ell \nu_\ell \) or \( D \rightarrow K \ell \nu_\ell \) decay processes assuming a polar \( t \)-dependence of the vector form factor as well as a determination of the weak current of the chiral theory [1] based on QCD sum rules [5]:

\[ L^\mu = i \frac{\alpha}{2} < \gamma^\mu (1 - \gamma_5) H_b \xi^\dagger \xi_b >. \]

The parameter \( \alpha \) is related to the leptonic constant \( f_P \) defined by:

\[ < 0| \bar{q}^a \gamma^\mu \gamma_5 Q |P_a(p) > = if_P p^\mu, \]

by the relation \( \alpha = f_P \sqrt{m_P} \) (neglecting logarithmic corrections). Since \( 1/m_Q \) corrections appear to be significant in the \( f_B/f_D \) ratio, one could choose to include them and determine \( g \) from the \( D \rightarrow \pi, K \ell \nu_\ell \) decays by using for \( f_D \) the value \( \simeq 200 \, MeV \), indicated by both QCD sum rules [3] and lattice calculations [6]. However, as shown in [7], this
procedure produces weak hadronic matrix elements which, when used to predict non leptonic $B$ decays in the factorization approximation, give results in disagreement with data. Therefore, following \[7\] we choose to fix $\alpha$ from the value of $f_B$ obtained by QCD calculations \[5, 8\] which, given the actual value of $m_b$, much larger than $m_c$, should provide a better approximation. In \[8\] the QCD sum rules method is applied to the evaluation of $f_B$ in the infinite heavy quark mass limit, with the result:

$$\alpha = 0.35 - 0.45 \text{ GeV}^{3/2}. \quad (16)$$

Theoretical uncertainties of this result are significant because the $O(\alpha_s)$ corrections are more than 50% of the result. Nevertheless the result \((16)\) is confirmed by another \[8\] QCD sum rules calculation of $f_B$, performed for finite value of $m_b$, which gives $\alpha = f_B \sqrt{M_B} \simeq 0.47 \pm 0.04 \text{ GeV}^{3/2}$. Using \((16)\) and semileptonic $D$ decays data ($D \to \pi \ell \nu$ and $D \to K \ell \nu$) we get:

$$|g| = 0.40 \pm 0.08, \quad (17)$$

where the error is only theoretical and is a consequence of the uncertainty in eq.\((16)\); to this theoretical error one should add a further 20% experimental uncertainty from $D \to \pi(K) \ell \nu$ decay data which give $|\lambda \alpha| = 0.16 \text{ GeV}^{3/2}$. Using the same value for $\alpha$, from $D \to K^* \ell \nu_{\ell} \nu$ decay data \[8\] one obtains for $\lambda$:

$$|\lambda| = 0.40 \pm 0.08 \text{ GeV}^{-1}, \quad (18)$$

where again the error is only theoretical (the experimental error is a further 20%).

Let us now turn to the decay:

$$D^* \to D \gamma, \quad (19)$$

whose matrix element can be written as follows:

$$\mathcal{M}(D^* \to D \gamma) = e e^{*\mu} J_{\mu} \quad (20)$$

with:

$$4$$
\[ J_\mu = \langle D(p')|J^{em}_\mu|D^*(p, \eta) \rangle = \langle D(p')|e_Q \bar{Q} \gamma^\mu Q + e_q \bar{q} \gamma^\mu q|D^*(p, \eta) \rangle = e_Q J^Q_\mu + e_q J^q_\mu, \] (21)

where \( e_Q = \frac{2}{3} \) is the heavy quark \((Q = c)\) charge and \( e_q \) is the light quark charge \((e_q = e_u = 2/3\) for \( D^{*0} \) and \( e_q = -1/3 \) for \( D^{*+} \) and \( D^*_s \)). Let us consider the two currents appearing in (21) separately. \( J^Q_\mu \) can be expressed in terms of the Isgur-Wise universal form factor \([10]\) as follows:

\[ \langle D(p')|\bar{c} \gamma^\mu c|D^*(p, \eta) \rangle = i \sqrt{m_D m_{D^*}} \xi(v \cdot v') \epsilon^{\mu \nu \alpha \beta} \eta_{\nu} v_{\alpha} v'_{\beta}, \] (22)

where \( p' = m_D v', p = m_{D^*} v \) and \( v \cdot v' \approx 1 \) because:

\[ 0 = q^2 = m_D^2 + m_{D^*}^2 - 2 m_D m_{D^*} v \cdot v'. \] (23)

As to the computation of the current

\[ J^q_\mu = \langle D(p')|\bar{q} \gamma^\mu q|D^*(p, \eta) \rangle, \] (24)

we assume VMD hypothesis and write:

\[ J^q_\mu = \sum_{V, \lambda} \frac{i}{q^2 - m_V^2} < D(p')V(q, \epsilon_1(\lambda))|D^*(p, \eta) \rangle \langle 0|\bar{q} \gamma^\mu q|V(q, \epsilon_1(\lambda)) \rangle, \] (25)

where \( q^2 = 0 \) and the sum is over the vector meson resonances \( V = \omega, \rho^0, \phi \) and over the \( V \) helicities. The vacuum-to-meson current matrix element appearing in (25) is given, assuming \( SU(3) \) flavour symmetry, by:

\[ < 0|\bar{q} T^i \gamma^\mu q|V(q, \epsilon_1) \rangle = \epsilon_1^\mu f_V Tr(V T^i), \] (26)

where \( (T^i)_{mn} = \delta_{im} \delta_m \) and \( i = 1, 2, 3 \) for \( u, d, s \) respectively. From \( \omega \rightarrow e^+ e^- \) and \( \rho^0 \rightarrow e^+ e^- \) decays \([11]\) we get the same value for \( f_V \): \( f_V = 0.17 \text{ GeV}^2 \); from \( \phi \rightarrow e^+ e^- \) we
have \( f_\phi = f_V + \delta f \) with \( \delta f = 0.08 GeV^2 \), showing a significant \( SU(3) \) violation. Using (26) and the strong lagrangian \( L_2 \) we can easily compute \( J_\mu^a \) and therefore (21). The results are:

\[
\mathcal{M}(D^* \rightarrow D \gamma) = i \epsilon^{\mu\nu\alpha\beta} \eta^\nu v_\alpha v'_\beta \sqrt{m_{D^*} m_D} \left[ e_Q - e_q \frac{2\sqrt{2} g_V \lambda m_{D^*} f_V}{m_\omega^2} \right],
\]

\[
\mathcal{M}(D^*_s \rightarrow D_s \gamma) = i \epsilon^{\mu\nu\alpha\beta} \eta^\nu v_\alpha v'_\beta \sqrt{m_{D^*_s} m_{D^*_s}} \left[ e_Q + \frac{1}{3} 2\sqrt{2} g_V \lambda m_{D^*_s} f_{\phi} \right],
\]

where \( e_Q = e_c = \frac{2}{3} \). Eq. (27) holds for both \( D^{*+} \rightarrow D^{+} \gamma \) and \( D^{*0} \rightarrow D^{0} \gamma \) (with \( e_q = -\frac{1}{3} \) and \( \frac{2}{3} \) respectively), assuming \( m_{\rho}^2 \approx m_{\omega}^2 \). Since (18) only gives the absolute value of \( \lambda \), we have fixed its sign by imposing that the relative sign between the two contributions is identical to the one given by the constituent quark model [12], i.e. we take \( \lambda = -0.40 \pm 0.08 \). It is worth observing that eqs. (27, 28) describe with obvious changes also \( B^* \) radiative decays.

From the amplitudes (27, 28) and from eq. (13) we can compute decay rates and branching ratios (BR) for \( D^* \) and \( B^* \) decays. They are reported in Table I together with the CLEO data [13] on radiative \( D^* \) decays. We observe an overall agreement between theoretical results and experiment. In particular, the tiny decay rate \( D^{*+} \rightarrow D^{+} \gamma \) can be explained as an effect of a cancellation between the two contributions appearing in (27).

Let us compare our results with previous work. The general structure of the matrix elements (27) and (28) coincides with previous analyses [12, 14, 15, 16, 17]; the main differences are in the determination of the light quark current that is not provided by the heavy quark effective theory. Our outcome coincides with the constituent quark model result: \( \mathcal{M}(D^* \rightarrow D \gamma) \sim \left( \frac{e_Q}{m_Q} + \frac{e_q}{m_q} \right) \) with \( (m_q)^{-1} = -2\sqrt{2} \lambda g_V \frac{f_{\omega}}{m_V^2} \) (see [12, 17]). Whereas \( m_q \) in the quark model coincides with the constituent light quark mass (e.g. \( m_u \approx m_d \approx 300 MeV \) [12], in our case the mass parameter \( -2\sqrt{2} \lambda g_V f_V / m_\omega^2 \) has a significantly larger value, i.e. 0.55 GeV and 0.95 GeV for \( D^* \) and \( D^*_s \) decays respectively.

The input value for \( g \) (eq. (17)) and the result for \( (m_q)^{-1} \) are compatible with the analyses of ref. [13] where \( g \) and \( m_q \) are treated as free parameters; in that paper also \( SU(3) \) breaking through one loop diagrams are reported. The breaking of \( SU(3) \) given by (28) is actually different from that one, since it is due to explicit \( SU(3) \) violation (by \( m_\phi \neq m_\omega \) and \( f_\phi \neq f_\omega \)) and not to chiral loop effects.
Let us now turn to the radiative decays of positive parity charmed meson resonances. As well known \[21, 22\] in HQET one expects four states: two of these states, the \( J_P = 0^+ \) (\( D_0 \)) and \( J_P = 1^+ \) (\( D_1^0 \)) form a mass doublet and are characterized by a total angular momentum of the light degrees of freedom \( s_\ell = \frac{1}{2} \); a mass doublet is also formed by the two states having \( s_\ell = \frac{3}{2} \) with \( J_P = 1^+ \) (\( D_1^+ \)) and \( J_P = 2^+ \) (\( D_2^+ \)) \[3\]. These states are described by effective operators analogous to the fields \( H_a \) introduced in eq.(7); in particular, the states having \( s_\ell = \frac{1}{2} \) are described by:

\[
S_a = \frac{1 + \gamma_5}{2} [D_1^{\mu} \gamma_\mu \gamma_5 - D_0] ;
\]

for these states we write the effective couplings to the light vector and heavy negative parity mesons as follows \[2\]:

\[
L' = i\zeta < \bar{S}_a H_\rho \gamma_\mu (\mathcal{V}_\mu - \rho_\mu)_{ba} > + i\mu < \bar{S}_a H_b \sigma^{\lambda\nu} F_{\lambda\nu} (\rho)_{ba} > + \text{h.c.}
\]

The analysis of the \( D \to K^* \) semileptonic decays via axial current, using a polar dependence of the form factors \[4\], provides the result \[2\]:

\[
\mu = -0.13 \pm 0.05 \text{ GeV}^{-1}
\]

Similarly to the \( D^* \) radiative decays, in order to describe the radiative processes

\[
\begin{align*}
D_0 & \to D^* \gamma \\
D_1' & \to D \gamma \\
D_1' & \to D^* \gamma
\end{align*}
\]

we distinguish in the electromagnetic matrix element the term containing the heavy quark current \( e_c \bar{c} \gamma_\mu c \) and the term with the light quarks \( e_q \bar{q} \gamma_\mu q \). The first can be obtained from the Lagrange density:

\[
\mathcal{L}'' = \frac{e}{2m_Q} \epsilon_Q \bar{h}_\nu \sigma^{\mu\nu} h_v F_{\mu\nu} ,
\]
which allows the transition $Q \rightarrow Q \gamma$ \footnote{The use of (33) for $D^* \rightarrow D \gamma$ produces exactly the same result already obtained in the leading $1/m_Q$ expansion.}; the matrix elements involve the universal form factor $\tau_{1/2}$, analogous to the Isgur-Wise function, which has been introduced in \cite{21} and computed in \cite{24} by QCD sum rules. On the other hand, the matrix element of the light quark current $e_q\bar{q}\gamma_\mu q$ can be related, via VMD, to the coupling in eq. (30). However, in this case the contribution of the $\zeta$ term, used together with the VMD hypothesis, displays a breaking of gauge invariance due to non-leading $1/m_Q$ terms. In principle, gauge invariance could be recovered adding further couplings, but this procedure would spoil the predictivity of the method. Since no experimental information is available on $\zeta$, we choose $\zeta = 0$ in our estimate of the positive parity radiative transitions. This should give at least an order of magnitude estimate of the $0^+, 1^+$ radiative decay widths. The obtained results are given in Table II. We observe that the radiative widths of positive resonances are small due to an almost complete cancellation between the two contributions of the e.m. current. We also note that for the neutral charged resonances, the values in Table II would correspond to branching ratios of the order $10^{-4} - 10^{-3}$ \cite{22, 23}.

In conclusion, we have shown that chiral heavy meson theory can be used to relate semileptonic $B$ and $D$ decays to charmed resonances decays by using the additional hypothesis of Vector Meson Dominance for the light quark vector current. Even though our results have potentially sizable uncertainties ($1/m_Q$ corrections to HQET, correction to VMD, etc.), by using semileptonic decays as an input we have found results that are in fairly good agreement with the $D^*$ decay rates; moreover, heavy flavour and chiral symmetries have been used to predict the decay width of $D_s^*$ and $B^*$ mesons.
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Tables Captions

Table I Theoretical and experimental $D^*$ and $B^*$ decay rates.

Table II Radiative decay widths of positive parity charmed mesons.
### Table I

| Decay rate/ BR | theory | experiment |
|----------------|--------|------------|
| $\Gamma(D^{*+})$ | $46.1 \pm 14.2$ KeV | $< 131$ KeV [23] |
| $BR(D^{*+} \to D^{+}\pi^0)$ | $31.2 \pm 17.4\%$ | $30.8 \pm 0.4 \pm 0.8$ |
| $BR(D^{*+} \to D^0\pi^+)$ | $67.7 \pm 34.2\%$ | $68.1 \pm 1.0 \pm 1.3$ |
| $BR(D^{*+} \to D^{+}\gamma)$ | $1.1 \pm 0.9\%$ | $1.1 \pm 1.4 \pm 1.6$ |

| $\Gamma(D^{*0})$ | $36.7 \pm 9.7$ KeV |
| $BR(D^{*0} \to D^{0}\pi^0)$ | $56.4 \pm 27.1\%$ | $63.6 \pm 2.3 \pm 3.3$ |
| $BR(D^{*0} \to D^{0}\gamma)$ | $43.6 \pm 17.8\%$ | $36.4 \pm 2.3 \pm 3.3$ |

| $\Gamma(D^*_s) = \Gamma(D_s^* \to D_s\gamma)$ | $(0.24 \pm 0.24)$ KeV |

| $\Gamma(B^{*+}) = \Gamma(B^{*+} \to B^{+}\gamma)$ | $(0.22 \pm 0.09)$ KeV |

| $\Gamma(B^{*0}) = \Gamma(B^{*0} \to B^{0}\gamma)$ | $(0.075 \pm 0.027)$ KeV |

### Table II

| Decay mode | width (KeV) |
|------------|-------------|
| $\Gamma(D^{*0}_1 \to D^{*0}\gamma)$ | $93 \pm 44$ |
| $\Gamma(D^{*0}_1 \to D^0\gamma)$ | $14 \pm 6$ |
| $\Gamma(D^{*0}_0 \to D^{*0}\gamma)$ | $115 \pm 54$ |
| $\Gamma(D^{*+}_1 \to D^{*+}\gamma)$ | $< 2.3$ |
| $\Gamma(D^{*+}_1 \to D^{+}\gamma)$ | $< 3.3$ |
| $\Gamma(D^{*+}_0 \to D^{+}\gamma)$ | $< 2.8$ |