Abstract

Electromagnetic interactions are introduced in the effective chiral Lagrangian for heavy mesons which includes light vector particles. A suitable notion of vector meson dominance is formulated. It is found that the heavy meson - light vector and heavy meson - light pseudoscalar coupling constants obtained from $D^- \to \pi^- \gamma$ branching ratios are compatible with those estimated from semi-leptonic transition amplitudes. It is found that the heavy meson - light vector meson dominance is formulated. The coupling constants for heavy mesons which include light vector particles are introduced in the effective chiral Lagrangian.
I. INTRODUCTION

Effective Lagrangians combining heavy quark symmetry and chiral invariance [1] provide promising tools for understanding the “soft” interactions of the heavy mesons. The apparent dominance of the decays \( D \to K^* l\nu \) over \( D \to K\pi l\nu \) [2] as well as general considerations [1] suggest the inclusion of the light vector mesons in addition to the light pseudoscalars. The total Lagrangian is the sum of a “light” part describing the three flavors \( u, d, s \) and a “heavy” part describing the “heavy” meson multiplet \( H \) and its interaction with the light sector:

\[
\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{light}} + \mathcal{L}_{\text{heavy}}.
\]  

The relevant light fields belong to the \( 3 \times 3 \) matrix of pseudoscalars \( \phi \Gamma \) and to the \( 3 \times 3 \) matrix of vectors \( \rho_\mu \). It is convenient to define objects which transform simply under the action of the chiral group \( \Gamma \)

\[
\xi = \exp \left( \frac{i \phi}{F_\pi} \right), \quad U = \xi^2,
\]

\[
A^L_\mu = \xi \rho_\mu \xi^\dagger + \frac{i}{\tilde{g}} \xi \partial_\mu \xi^\dagger,
\]

\[
A^R_\mu = \xi^\dagger \rho_\mu \xi + \frac{i}{\tilde{g}} \xi^\dagger \partial_\mu \xi,
\]

\[
F_{\mu\nu} = \partial_\mu \rho_\nu - \partial_\nu \rho_\mu - i \tilde{g} [\rho_\mu, \rho_\nu],
\]

where \( F_\pi \approx 0.132 \text{ GeV} \) and \( \tilde{g} \approx 3.93 \) for a typical fit.

The heavy multiplet field combining the heavy pseudoscalar \( P' \) and the heavy vector \( Q'_\mu \Gamma \) both moving with a fixed 4-velocity \( V_\mu \) is given by

\[
H = \frac{1 - i \gamma_\mu V_\mu}{2} (i \gamma_5 P' + i \gamma_\nu Q'_\nu), \quad \bar{H} = \gamma_4 H\gamma_4.
\]  

In our convention \( H \) has the canonical dimension one.

The light part of the action under consideration has been most recently discussed in [3]. Apart from \( \text{SU}(3) \) and chiral symmetry breaking terms and terms proportional to the Levi-Civita symbol \( \Gamma \) it may be written as

\[
\mathcal{L}_{\text{light}} = -\frac{1}{4} Tr \left( F_{\mu\nu}(\rho) F_{\mu\nu}(\rho) \right) - \frac{m_v^2 (1 + k)}{8k} Tr \left( A^L_\mu A^L_\mu + A^R_\mu A^R_\mu \right) + \frac{m_v^2 (1 - k)}{4k} Tr \left( A^L_\mu U A^R_\mu U^\dagger \right),
\]

where \( m_v \approx 0.77 \text{ GeV} \) is the light vector mass and \( k = \frac{m_v^2}{F_\pi \tilde{g}} \). An alternate “hidden symmetry” approach [4] leads to the identical Lagrangian.

\( \mathcal{L}_{\text{heavy}} \) has been discussed by several authors [5–8]. Following the notations of ref. [6] it is

\[
\frac{\mathcal{L}_{\text{heavy}}}{M} = i V_\mu Tr \left[ H (\partial_\mu - i \tilde{g} \rho_\mu - i (1 - \alpha) v_\mu) \bar{H} \right] + i d Tr \left[ H \gamma_\mu \gamma_5 p_\mu \bar{H} \right] + \frac{i c}{m_v} Tr \left[ H \gamma_\mu \gamma_\nu F_{\mu\nu}(\rho) \bar{H} \right],
\]  

(1.5)
where $M$ is the mass of the heavy meson and

$$v_\mu, p_\mu = \frac{i}{2} (\xi \partial_\mu \xi^\dagger \pm \xi^\dagger \partial_\mu \xi). \quad (1.6)$$

$\alpha, c, d$ are dimensionless coupling constants for the heavy-light interactions; they are crucial for discussing the soft dynamics of the heavy mesons as well as other applications. At present we are at a rather preliminary stage on the route to their precise determination. The main purpose of this note is to estimate the parameter $c$ from experimental information on the $D^* \rightarrow D \gamma$ rates based on a suitable notion of vector meson dominance for the electromagnetic interactions of heavy mesons. Analogous considerations have been presented for the model without light vector mesons [9-11]. At the present stage it seems most reasonable to work in the leading order in $M$ as well as tree order in the light fields.

The coupling constant $d$ in (1.5) is related to the $D^* \rightarrow D \pi$ decay widths as follows:

$$
\begin{align*}
\Gamma(D^{*0} \rightarrow D^0 \pi^0) &= \frac{d^2 p_0^2}{12\pi F_\pi^2}, \\
\Gamma(D^{*+} \rightarrow D^0 \pi^+) &= \frac{d^2 p_0^2}{6\pi F_\pi^2}, \\
\Gamma(D^{*+} \rightarrow D^+ \pi^0) &= \frac{d^2 p_0^2}{12\pi F_\pi^2}.
\end{align*}
$$

where $p_\pi$ is the decay 3-momentum in the parent rest frame for the $D^* \pi^0$ final state. Note that $D^{*0} \rightarrow D^0 \pi^0$ in addition to all of the $B^* \rightarrow B \pi$ decays are energetically forbidden. $D^{*+} \rightarrow D^0 \pi^0$ is energetically allowed but is suppressed [12] due to isospin conservation and has not yet been observed. Since the total width of $D^*$ is not known to the relevant accuracy no useful statement can be made about the constant $d$ from (1.7) alone. Rather (1.7) must be used together with information about the $D^* \rightarrow D \gamma$ decays.

The coupling constant $\alpha$ in (1.5) was introduced in [6] as a measure of vector meson dominance in the light-heavy direct interaction. We will also verify here that $\alpha = 1$ corresponds to vector meson dominance for the diagonal matrix elements of the light electromagnetic current between heavy meson states $\Gamma(A|J^\mu_{light}|A)\Gamma$ where $A = P$ or $Q_\mu$.

II. HEAVY VECTOR MESON RADIATIVE DECAYS

The decays of the type $D^* \rightarrow D \gamma$ are governed by the fundamental electromagnetic interaction:

$$\mathcal{L}_{EM} = eJ^EM_\mu A_\mu \quad (2.1)$$

where $e$ is the proton charge and $A_\mu$ is the photon field. It is important to note that we need both the light and the heavy pieces in the decomposition:

$$J^EM_\mu = J^light_\mu + J^{heavy}_\mu,$$

$$J^light_\mu = \frac{i}{3} \bar{u} \gamma_\mu u - \frac{1}{3} \left( \bar{d} \gamma_\mu d + \bar{s} \gamma_\mu s \right),$$

$$J^{heavy}_\mu = iC\bar{Q} \gamma_\mu Q + \cdots, \quad (2.2)$$
where \( C \) is the electric charge of the particular heavy quark under consideration (e.g. \( \frac{2}{3} \) for the \( c \) quark).

For orientation purposes it is useful to consider how the \( D^* \to D \gamma \) decays are computed in the simplest non-relativistic “constituent” quark model [13]. The interaction Hamiltonian is \( -\mu \cdot B \) where the magnetic moment operator is

\[
\mu = e \left( \frac{C}{M} S^{\text{heavy}} + \frac{\bar{q}}{m} S^{\text{light}} \right),
\]

while \( M \) and \( m \) are the heavy and light constituent quark masses respectively. \( \bar{q} \) is the charge of the the light anti-quark in the heavy meson. The two terms in (2.3) illustrate the decomposition of \( J^E_M \) in (2.2). After sandwiching (2.3) between vector and pseudoscalar spin wavefunctions we find that the amplitude for \( D^* \to D \gamma \) is proportional to \( (\frac{C}{M} - \frac{\bar{q}}{m}) \).

This means that the amplitudes for \( D^{*0} \to D^0 \gamma \), \( D^{*+} \to D^+ \gamma \), and \( D^{*0} \to D^+ \gamma \) stand in the ratios:

\[
(1 + \frac{m}{M}) : (-\frac{1}{2} + \frac{m}{M}) : (-\frac{m}{2m_s} + \frac{m}{M})
\]

where \( m_s \) is the constituent strange quark mass. (For comparison the corresponding radiative amplitudes for \( \bar{B}^* \to \bar{B}^- \gamma \), \( \bar{B}^{*0} \to \bar{B}^0 \gamma \) and \( \bar{B}^{*0} \to \bar{B}^0 \gamma \) stand in the ratios \( (-2 + \frac{m}{M}) : (\frac{m}{m_s} + \frac{m}{M}) \)). Now we would like to start working in the leading \( M \to \infty \) limit. However \( \Gamma \) indicates that this is a rather questionable approximation for \( D^* \to D \gamma \) since \( m \approx 0.35 \text{ GeV} \) while \( M \) is in the 1.5 - 1.8 GeV range. In the case of \( D^{*+} \to D^+ \gamma \) the piece proportional to \( \frac{1}{M} \) is expected to be almost half of the leading term and opposite in sign.

A possible approach in the more general case is to include the \( \frac{1}{M} \) corrections corresponding to \( J^H_M \) in (2.2) while neglecting the \( \frac{1}{M} \) corrections to \( J^L_M \). This may be reasonable since it has been pointed out [10] that the contribution of \( J^H_M \) to the radiative decays is of the order \( \frac{1}{M} \) and is fixed from its relation to the heavy quark number current. This corresponds to a term in the effective Lagrangian:

\[
-\frac{1}{2} e C A_{\mu} Tr[\partial_\mu (H \tilde{R} \sigma_{\mu\nu})].
\]

### III. VECTOR MESON DOMINANCE

First let us review how to add electromagnetic interactions to the light particle Lagrangian (1.4). More details are given in [14] and in section III of [15]. The fields \( A^L_\mu \) and \( A^R_\mu \) introduced in (1.2) are taken to transform under a local chiral transformation \( U_{L,R} = 1 + E_{L,R} \) as

\[
\delta A^{L,R}_\mu = -[A^{L,R}_\mu, E_{L,R}] - \frac{i}{\bar{g}} \partial_\mu E_{L,R}.
\]

\(^1\bar{g} \) is called \( g \) in these references.
External fields $B^{L,R}_\mu$ transform as
\[ \delta B^{L,R}_\mu = - [B^{L,R}_\mu, E_{L,R}] - \frac{i}{\hbar} \partial_\mu E_{L,R}, \] (3.2)
where $\hbar$ is the external field coupling constant. Then it is clear that (1.4) will become locally gauge invariant if we make the substitutions $A^{L,R}_\mu \rightarrow A^{L,R}_\mu - \frac{\hbar}{2} B^{L,R}_\mu$. Here we are interested in the case of electromagnetism which corresponds to the choice
\[ h B^{L,R}_\mu = e Q A_\mu, \quad Q = \text{diag}(\frac{2}{3}, -\frac{1}{3}, -\frac{1}{3}). \] (3.3)
The resulting electromagnetic interaction piece from (1.4) may be expanded out (see (14) of [14]) to yield the leading relevant terms
\[ e A_\mu [\bar{q} g F^2_x T_r(Q \rho_\mu) + i (1 - \frac{k}{2}) T_r(Q(\bar{\phi} \partial_\mu \phi - \partial_\mu \phi \bar{\phi})) + \cdots]. \] (3.4)
It is possible to eliminate the photon-vector meson cross term by rediagonalization but it is more conventional to keep it in this form. Notice that the special choice $k = 2$ is denoted the Kawarabashi-Suzuki-Riazuddin-Fayazuddin (KSRF) relation [16]. When this holds the second term in (3.4) vanishes and the photon couples to the charged light pseudoscalars via its mixing with the light neutral vector mesons in the first term. Actually $k$ seems to be 10% higher than the value required by the KSRF relation so that this picture is reasonably good but not perfect.

A leading $O(M_0^0)$ contribution to the $D^* \rightarrow D \gamma$ decays arises via the $c$ term in (1.5) (which is locally chiral invariant) giving $D^* D \rho^0$ and $D^* D \omega^0$ vertices followed by the $\rho^0 - \gamma$ and $\omega^0 - \gamma$ mixings in (3.4). In addition it is possible to construct a direct $H \bar{H} \gamma$ locally chiral invariant interaction term as:
\[ \frac{i e M_\delta}{2 m_v} T_r(H^n_{\mu \nu} \xi^\dagger F_{\mu \nu}(B_L) \xi + \xi F_{\mu \nu}(B_R) \xi^\dagger) \bar{H}, \] (3.5)
where we used the facts that under local chiral transformations $\xi \bar{H} \rightarrow U_L(x) \xi \bar{H}$ and $\xi^\dagger \bar{H} \rightarrow U_R(x) \xi^\dagger \bar{H}$. Specializing this to electromagnetic external fields using (3.3) gives an effective term:
\[ \frac{i e M_\delta}{2 m_v} F_{\mu \nu}(A) T_r(H^n_{\mu \nu} \gamma_\nu [\xi^\dagger Q \xi + \xi Q \xi^\dagger] \bar{H}), \] (3.6)
whose strength is measured by the parameter $\delta$.

Putting the contributions to $D^* \rightarrow D \gamma$ from [(1.5) and (3.4)] and from (3.6) together with the subleading one from (2.5) finally yields the width expressions:
\[ \Gamma(D^0 \rightarrow D^\gamma) = \frac{e^2 p_0^3}{12 \pi} \left[ \frac{2}{3 M} + \frac{8}{3 m_v} (\delta + \frac{c}{g}) \right]^2, \]
\[ \Gamma(D^{*+} \rightarrow D^{+} \gamma) = \frac{e^2 p_\gamma^3}{12 \pi} \left[ \frac{2}{3 M} - \frac{4}{3 m_v} (\delta + \frac{c}{g}) \right]^2, \]
\[ \Gamma(D^{*+} \rightarrow D^{+*} \gamma) \approx \Gamma(D^{*+} \rightarrow D^{+} \gamma), \] (3.7)
wherein $p_0$, and $p_{\gamma}$ are the 3-momenta in the $D^{*0}$ and $D^{*+}$ rest frames respectively. The third approximate equality utilizes the coincidence that the phase space factors are approximately equal for the two reactions. We shall not make use of the $D^{*+} \rightarrow D^+ \gamma$ reaction here. To compute it more accurately in the present model (even at the tree level) involves taking into account several SU(3) symmetry breaking terms discussed in [3]. Notice that the structure of the amplitudes in (3.7) is essentially the same as that in the naive quark model[2.4]. The parameter $M$ in (3.7) is however more reasonably considered to be the heavy meson mass. A natural notion of light vector meson dominance for $D^* \rightarrow D \gamma$ is to set

$$\delta = 0$$ (3.8)

which corresponds to the photon interacting with the light electromagnetic transition moment only through its mixing with the light vectors in (3.4). Of course there is no a priori reason for the assumption (3.8) to be perfect but usual low energy phenomenology suggests that it is very sensible.

Now let us compare with experiment. The latest data from CLEO II [17] yields

$$\frac{\Gamma(D^{*0} \rightarrow D^0 \gamma)}{\Gamma(D^{*0} \rightarrow D^0 \pi^0)} = 0.57 \pm 0.14,$$

$$\frac{\Gamma(D^{*+} \rightarrow D^+ \gamma)}{\Gamma(D^{*+} \rightarrow D^0 \pi^+)} < 0.059.$$ (3.9)

These numbers should be equated to the predictions from (3.7) and (1.7). Defining for temporary convenience

$$A = \frac{2}{3M_d}, \quad B = \frac{4}{3m_v d}(\delta + \frac{c}{\bar{g}}),$$ (3.10)

we then get $|A + 2B| = 3.40 \pm 0.42 \text{ GeV}^{-1}$ as well as $|A - B| < 1.37 \text{ GeV}^{-1}$. We should bear in mind that none of the four quantities $M, d, c, \delta$ involved in (3.10) are really known accurately. We first eliminate $A$ by assuming that the relative signs of $A$ and $B$ in (3.7) are the same as the corresponding $O(\frac{1}{M})$ and $O(M^0)$ terms in the quark model amplitudes (2.4). This results in

$$1.13 \pm 0.14 < |B| < 1.59 \pm 0.14 \text{ GeV}^{-1}.$$ (3.11)

If we now invoke the vector meson dominance assumption that $\delta = 0$ we are left with a fairly stringent bound on the interesting $|\frac{c}{d}|$ ratio:

$$2.56 \pm 0.32 < |\frac{c}{d}| < 3.61 \pm 0.32.$$ (3.12)

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3We would like to thank S. Stone for pointing out to us that the value quoted in [17] for $\frac{\Gamma(D^{*+} \rightarrow D^+ \gamma)}{\Gamma_{\text{tot.}}}$ is best interpreted as an upper bound of 4%
As we will discuss in the next section, this appears to be very reasonable. Possible conflict with other determinations of \( c \) and \( d \) and improved experimental accuracy might be resolved by relaxing the \( \delta = 0 \) assumption.

Another option for our analysis is to eliminate \( B \) instead of \( A \). Then taking both \( A \) and \( B \) positive, we easily see that as \( (B-A) \) ranges from 0 to 1.37:

\[
0.59 \pm 0.07 \text{GeV} < M_d \leq 3.05 \pm 1.94 \text{GeV}.
\] (3.13)

Taking \( M=1.8 \text{ GeV} \) for example, leads to \( d \) in the approximate range 0.28 to 2.77.

To end this section we comment on the first kinetic term in (1.5) which contains the “chiral derivative” \( \mathcal{D}_\mu \bar{H} = [\partial_\mu - i \tilde{\alpha} \tilde{\rho}_\mu - i(1-\alpha)\nu_\mu] \bar{H} \). The presence of \( \mathcal{D}_\mu \) guarantees invariance under \textit{global} transformations belonging to the chiral group. When including electromagnetism we should naturally add the terms \(-ieA_\mu (Q-C)\) to \( \mathcal{D}_\mu \). This is not sufficient since both \( \nu_\mu \) and \( \rho_\mu \) pick up inhomogenous pieces under \textit{local} electromagnetic U(1) transformations. Thus we should replace them by the properly covariant quantities

\[
\bar{\nu}_\mu = \nu_\mu + eA_\mu \frac{1}{2} (\xi Q \xi^\dagger + \xi^\dagger Q \xi) - 2Q],
\]

\[
\bar{\rho}_\mu = \rho_\mu - \frac{e}{g} Q A_\mu.
\]

The net result is an “electrified chiral derivative” to be used in (1.5):

\[
\mathcal{D}_\mu^{\text{EC}D} \bar{H} = [\partial_\mu - ieA_\mu (Q-C) - i \tilde{\alpha} \tilde{\rho}_\mu - i(1-\alpha)\bar{\nu}_\mu] \bar{H}
= \mathcal{D}_\mu \bar{H} + ieA_\mu [C - \frac{1}{2}(1-\alpha)(\xi Q \xi^\dagger + \xi^\dagger Q \xi)] \bar{H}.
\] (3.14)

From this expression, it is evident that the choice \( \alpha = 1 \) corresponds to no direct photon coupling to the \textit{light} part of the heavy meson field \( \bar{H} \). The indirect coupling via photon mixing with the light vectors ensures proper normalization of the electromagnetic form factors of the heavy meson at zero momentum transfer. Since the first term in (1.5) is diagonal, of course, it does not contribute to the off-diagonal transition matrix elements of interest. Even though \( \alpha = 1 \) has been seen to be the choice for vector meson dominance of the diagonal matrix elements of light electromagnetic currents between heavy meson states, one still does not know just how good that assumption really is.

\section{IV. OTHER ESTIMATES AND DISCUSSION}

In order to compare with the vector meson dominance estimate (3.12) we now consider some other estimates for the light pseudoscalar-heavy meson coupling constant \( d \) and the light vector-heavy meson coupling constant \( c \).

One way to get a handle on \( c \) and \( d \) is to imagine that the flavor SU(3) invariant expression for the vector-vector-pseudoscalar interaction [see (2.18) of [15] for example]:

\[
\mathcal{L}_{VV \phi} = -ig_{VV \phi} \epsilon_{\mu\nu\gamma\delta} Tr(\partial_\mu \rho_\nu \partial_\gamma \rho_\delta \phi),
\] (4.1)

continues to hold when the strange quark is formally considered to be “heavy”. Then the \( K^+ \) field, normally denoted as \( \phi_1 \), would be called \( \bar{P}_1 \) while the \( K^{*+} \) field would be called...
\( \bar{Q}_1 \mu \) (in the notation of [6]). If we consider both of the vectors in (4.1) to be heavy (with the pseudoscalar light) the resulting \( Q\phi\bar{Q} \) interaction is actually part of the \( d \)-term in (1.5) (see (3.20) of [6]). Similarly the \( P\rho\bar{Q} \) piece from (4.1) is part of the \( e \)-term in (1.5) (see (3.22) of [6]). These identifications give the estimates:

\[
-g_{V\phi} = \frac{2d}{F_{\pi}} = \frac{4c}{m_{\nu}}.
\]

(4.2)

We immediately see that the ratio:

\[
\frac{c}{d} = \frac{m_{\nu}}{2F_{\pi}} = 2.92
\]

(4.3)

is compatible with (3.12). A typical estimate \([15] \Gamma |g_{V\phi}F_{\pi}| \approx 1.8 \) yields the additional expectation that \(|d| \approx 0.9\).

A more direct experimental approach to finding \( d \) is based on fitting [18] the \( D \to K \) semi-leptonic transition form factor to \( \frac{B}{\gamma^2 + m^2(D^*)} \Gamma \) where \( R=\text{constant} \). In the present model\( R = \frac{ME_{\pi}}{F_{\pi}} \Gamma \) where \( M \) is the \( D \) mass. This leads to (see (5.2) of [19])

\[
|d| \approx 0.53,
\]

(4.4)

where \( F_D \approx 0.25 \text{ GeV} \) was taken [20]. A similar approach to finding \( c \) can be based on the study of the \( D \to K^* \) semi-leptonic transition vector type form factor. There is more complication in this case but one gets [21,22]:

\[
|c| \approx 1.6.
\]

(4.5)

We see that the ratio of (4.5) to (4.4) is also compatible with the vector meson dominance bounds in (3.12). The value of \( d \) given in (4.4) is within the range (3.13).

It may be helpful to finally give some predictions resulting from the typical parameter choices \( d = 0.53, c = 1.6, \delta = 0 \) and \( M = 1.8 \text{ GeV} \). Then the branching ratios in (3.9) turn out to be \( \Gamma(D^{*0} \to D^{0}\gamma) / \Gamma(D^{*0} \to D^0\pi^0) = 0.55 \) and \( \Gamma(D^{*+} \to D^{+}\gamma) / \Gamma(D^{*+} \to D^0\pi^+) = 0.01 \). The total widths are estimated as \( \Gamma_{\text{tot}}(D^{*0}) = 0.056 \text{ MeV} \) and \( \Gamma_{\text{tot}}(D^{*+}) = 0.081 \text{ MeV} \). With a heavy mass choice of 5.28 GeV the radiative B widths are estimated as \( \Gamma(B^{*-} \to B^-\gamma) = 0.43 \text{ KeV} \) and \( \Gamma(B^{*0} \to B^0\gamma) = 0.14 \text{ KeV} \).

We have discussed the introduction of the electromagnetic interaction in the framework of the effective chiral Lagrangian for heavy mesons which includes light vector mesons. A suitable notion of light vector meson dominance was formulated. Application was made to the radiative decays of the \( D^* \) mesons with the goal of determining the light heavy coupling constants \( c \) and \( d \). It was found that the acceptable range of values for the assumption of vector meson dominance was compatible with information extracted from semi-leptonic \( D \) decays. The structure of the radiative amplitudes had the same form as in the simple quark model and we made the additional assumption that the relative signs of the \( O(M^0) \) and \( O(\frac{1}{M}) \) pieces were the same as in this simple model. Apart from the necessity to include the \( \frac{1}{M} \) piece describing the heavy part of the electromagnetic current (which is “accidentally” enhanced for \( D^* \) radiative decays) we worked to leading order \( M^0 \). Furthermore, higher derivatives, loops and SU(3) symmetry breaking were neglected. Together with more accurate measurements of \( \Gamma(D^{*+} \to D^{+}\gamma) \) these additional corrections may in the future greatly clarify the situation. Deviations from precise vector meson dominance might then emerge.
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