Strong and Radiative $D^*$ Decays

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Abstract

We use the relativistic light–front quark model to show that both strong and radiative $D^*$ decays are in good agreement with the 1992 CLEO II results. In particular the coupling for $D^* \rightarrow D\pi$ is consistent with the experimental upper limit. The key point is the relativistic treatment of the quark spin.
1 Introduction

The $D^*$ decays has been under both theoretical and experimental investigation for quite a long time. The strong $D^*$ decays, $D^* \to D\pi$, can be described by the lowest–order (in external momentum) effective Lagangian

$$\mathcal{L} = \left(\frac{m_{D^*}g}{f_\pi}\right)\hat{D}^\dagger \tau \cdot \partial_\mu \pi \hat{D}^* \mu,$$  \hspace{1cm} (1)

where $\hat{D}$ and $\hat{D}^*$ are the $D$ and $D^*$ isodoublet fields, and $\tau$ is the isospin matrices. The decay width of $D^{*+} \to D^0 \pi^+$, for example, reads

$$\Gamma_{D^{*+} \to D^0 \pi^+} = \frac{g^2|p_\pi|^3}{6\pi f_\pi^2}.$$  \hspace{1cm} (2)

The radiative $D^*$ decays $D^* \to D\gamma$, on the other hand, can be described by the decay amplitude

$$M(D^* \to D\gamma) = e_G \epsilon_{\alpha\nu\sigma\lambda} \xi^\alpha \epsilon^\nu p_D^\sigma p_D^\lambda,$$  \hspace{1cm} (3)

where $e_G/2$ is the transition magnetic moment, and $\xi^\alpha$ and $\epsilon^\nu$ are the polarization vectors of the photon and the $D^*$ meson, respectively. The radiative decay width is

$$\Gamma = (e_G)^2|p_\gamma|^3 \frac{12\pi}{12\pi}.$$  \hspace{1cm} (4)

In many theoretical studies, people use the non–relativistic quark model to calculate the radiative $D^*$ decays while they use various modifications of the $SU(4)$ symmetry method for the strong $D^*$ decays. Obviously, the treatment of the strong and the radiative $D^*$ decays is not on the same footing.

However one can also use the quark model to calculate the strong decays. Through the PCAC relation one can relate the strong coupling $g$ to the axial–current form factor $A_0(0)$

$$g = A_0^{D^*}(0).$$  \hspace{1cm} (5)
Here $A_0$ is one of the four form factors used to describe the matrix element of a vector meson $V$ decaying to a pseudo–scalar meson $X$ by the quark transition $Q' \to Q$:

$$
\langle X(p_X)| \bar{Q} \gamma_\mu (1-\gamma_5) Q' |V(p_V, \epsilon) \rangle = \frac{2V(q^2)}{m_X + m_V} i\epsilon_{\mu \nu \alpha \beta} \epsilon^\nu p_X^\alpha p_V^\beta - 2m_V \frac{(e^* \cdot p_X)}{q^2} q^\mu A_0(q^2)
$$

$$
- \left[ (m_X + m_V) \epsilon^\nu A_1(q^2) - \frac{(e^* \cdot p_X)}{m_X + m_V} (p_X + p_V)^\mu A_2(q^2) - 2m_V \frac{(e^* \cdot p_X)}{q^2} q^\mu A_3(q^2) \right],
$$

where $\epsilon$ is the polarization vector of the vector meson, $q = p_V - p_X$, and

$$
A_3(q^2) = \frac{m_X + m_V}{2m_V} A_1(q^2) - \frac{m_V - m_X}{2m_V} A_2(q^2), \quad A_3(0) = A_0(0).
$$

In Eq. (5), $A^{DD^*}_0(0)$ is the $A_0(q^2 = 0)$ form factor for $D^* \to D$ induced by the quark transition $u \to d$. In the non–relativistic quark model, one obtains $g \simeq 1$ [10]. In a rather different model, the BSW model [11], one has [5]

$$
g = A^{DD^*}_0(0) = \int dx \ d^2k_\perp \phi^*_D(x, k_\perp) \phi_{D^*}(x, k_\perp),
$$

where $\phi(x, k_\perp)$ is the momentum wavefunction. Since the wavefunctions of the $D$ and $D^*$ are quite similar, this model also has $g \simeq 1$. In fact, in the heavy $c$–quark limit $\phi_{D^*}(x, k_\perp)$ and $\phi_D(x, k_\perp)$ should be exactly the same, due to the heavy quark symmetry, and hence in the BSW model, $g = 1$ in this limit. These results of $g$ are, however, larger than the upper limit of $g$ obtained from the ACCMOR measurement $\Gamma_{D^*} < 131\text{keV}$ [2]:

$$
g < 0.7.
$$

There have been some recent studies on the $D^*$ decays [3 4 5 6]. In ref. [6] the radiative $D^*$ decays are treated with the non–relativistic quark model but the strong coupling $g$ is related to the quark–level axial–current form factor $g_A^{ud} = \frac{3}{4}$; this is obtained from fitting the measured nucleon axial–current form factor $g_A^{nucleon} = 1.25 = \frac{5}{3}g_A^{ud}$. Thus the coupling $g$ is $g = g_A^{ud} = \frac{3}{4}$. In refs. [4 5] heavy–meson chiral–perturbation theory [12 13] is used and all the $D^*$ decays can be described by two parameters. These parameters are the couplings of the chiral Lagrangian
and can therefore be determined only by fitting the experimental data. In a recent study \[9\], the radiative decays are calculated using the vector–dominance hypothesis while the strong coupling $g$ is extracted from the data on the decay $D \to \pi e\bar{\nu}_e$ using heavy–meson chiral–symmetry relations. The central value obtained for $g$ is $g \approx 0.4$. In this study the uncertainties associated with assumptions on the $f_+$ form factor of $D \to \pi e\bar{\nu}_e$ and on the decay constant $f_D$, etc., could be large. A value of $g \sim 0.4$ is also obtained \[13\] in a numerical solution to another type of relativistic quark model \[14\].

The failure of the non–relativistic quark model and the BSW model in explaining $g < 0.7$ has its roots in the non–relativistic treatment of the quark spin in these models. In this letter, we present a calculation of the strong and radiative decays of $D^*$ using the relativistic light–front quark model \[15, 16, 17\]. In our calculation we are able to treat on the same basis both the strong and radiative decays. Due to the relativistic treatment one can see unambiguously how the strong coupling $g$ can be consistent with the limit in Eq. (9). Our calculated branching ratios for the strong and radiative decays are in good agreement with the 1992 CLEO II measurement \[3\].

2 The light–front relativistic quark model

The relativistic light–front quark model was developed quite a long time ago and there have been many applications \[15, 16, 17\]. Here we give a brief introduction.

A ground–state meson $V(\bar{Q}q)$ with spin $J$ on the light front can be described by the state vector

$$|V(P, J_3, J)\rangle = \int d^3p_1 d^3p_2 \delta(P - p_1 - p_2) \sum_{\lambda_1, \lambda_2} \Psi_{J, J_3}(P, p_1, p_2, \lambda_1, \lambda_2)|Q(\lambda_1, p_1) \bar{q}(\lambda_2, p_2)\rangle,$$  \hspace{1cm} (10)
In the light–front convention, the quark coordinates are given by

\[
p_1^+ = x_1 P^+ , \quad p_2^+ = x_2 P^+ , \quad x_1 + x_2 = 1 , \quad 0 \leq x_{1,2} \leq 1 ,
\]

\[
p_{1\perp} = x_1 P_{\perp} + k_{\perp} , \quad p_{2\perp} = x_2 P_{\perp} - k_{\perp} . \tag{11}
\]

The quantities \(x_{1,2}\) and \(k_{\perp}^2\) are invariant under the kinematic Lorentz transformations. Rotational invariance of the wave function for states with spin \(J\) and zero orbital angular momentum requires the wave function to have the form \([15, 16]\) (with \(x = x_1\))

\[
\Psi^{J,J_3}(P, p_1, p_2, \lambda_1, \lambda_2) = R^{J,J_3}(k_{\perp}, \lambda_1, \lambda_2) \phi(x, k_{\perp}) , \tag{12}
\]

where \(\phi(x, k_{\perp})\) is even in \(k_{\perp}\) and

\[
R^{J,J_3}(k_{\perp}, \lambda_1, \lambda_2) = \sum_{\lambda, \lambda'} \langle \lambda | R_M^+(k_{\perp}, m_Q) | \lambda \rangle \langle \lambda_2 | R_M^+(-k_{\perp}, m_{\bar{q}}) | \lambda' \rangle C^{J,J_3}(\frac{1}{2}, \lambda; \frac{1}{2}, \lambda') . \tag{13}
\]

In Eq. \((13)\), \(C^{J,J_3}(\frac{1}{2}, \lambda; \frac{1}{2}, \lambda')\) is the Clebsh–Gordan coefficient and the rotation \(R_M(k_{\perp}, m_i)\) on the quark spins is the Melosh rotation \([18]\):

\[
R_M(k_{\perp}, m_i) = \frac{m_i + x_i M_0 - i \sigma \cdot (n \times k_{\perp})}{\sqrt{(m_i + x_i M_0)^2 + k_{\perp}^2}} , \tag{14}
\]

where \(n = (0, 0, 1)\), \(\sigma\) is the Pauli spin matrix, and

\[
M_0^2 = \frac{m_Q^2 + k_{\perp}^2}{x_1} + \frac{m_{\bar{q}}^2 + k_{\perp}^2}{x_2} . \tag{15}
\]

The spin wave function \(R^{J,J_3}(k_{\perp}, \lambda_1, \lambda_2)\) in Eq. \((13)\) can also be written as

\[
R^{J,J_3}(k_{\perp}, \lambda_1, \lambda_2) = \chi_{\lambda_1}^\dagger R_M^+(k_{\perp}, m_Q) S^{J,J_3} R_M^+(-k_{\perp}, m_{\bar{q}}) \chi_{\lambda_2}
\]

\[
= \chi_{\lambda_1}^\dagger U^{J,J_3}_V \chi_{\lambda_2} , \tag{16}
\]

where \(S^{J,J_3}\) is defined by

\[
S^{J,J_3} = \sum_{\lambda, \lambda'} | \lambda \rangle \langle \lambda' | C^{J,J_3}(\frac{1}{2}, \lambda; \frac{1}{2}, \lambda') . \tag{17}
\]
For the pseudoscalar and vector mesons,
\begin{align}
S^{0,0} &= \frac{i\sigma_2}{\sqrt{2}}, \quad S^{1,\pm} = \frac{1 \pm \sigma_3}{2}, \quad S^{1,0} = \frac{\sigma_1}{\sqrt{2}}.
\end{align}

The explicit expressions of $U^{1,j_3}$ in Eq. (16) can be found in ref. [17].

The matrix element of a vector meson $V(Q'\bar{q})$ decaying to a pseudo–scalar meson $X(Q\bar{q})$ is
\begin{align}
\langle X(p_X)|\bar{Q}\Gamma Q'|V(p_V, J_3)\rangle &= \int dx \, d^2k_{\perp} \sum_{\lambda_1, \lambda_2} \Psi_{X}^{\lambda_0,0} \bar{u}_Q \Gamma u_{Q'} \Psi_{V}^{1,j_3} \\
&= \int dx \, d^2k_{\perp} \frac{\phi_X^* \phi_V}{x} Tr \left[ U_{X}^{0,0} U_{\Gamma} U_{V}^{1,j_3} \right], \quad (19)
\end{align}

where $U_{\Gamma}$ is defined by
\begin{align}
\bar{u}_i^Q \Gamma u_j^{Q'} = \chi_i^\dagger U_{\Gamma} \chi_j,
\end{align}

In Eq. (19), we choose $q^+ = 0$ so $q^2 = -q_{\perp}^2$. In contrast to Eq. (19), the matrix element in the BSW model is given by
\begin{align}
\langle X(p_X)|\bar{Q}\Gamma Q'|V(p_V, J_3)\rangle &= \int dx \, d^2k_{\perp} \frac{\phi_X^* \phi_V}{x} Tr \left[ S_{X}^{0,0} S^{1,j_3} \right]. \quad (21)
\end{align}

Eq. (19) is in fact expected to be valid only for “good” currents such as $\Gamma = \gamma^+, \gamma^+\gamma_5, \ldots$. There are contributions other than the one given in Eq. (19) if a current is not a “good” current [16].

Using the “good” currents $\Gamma = \gamma^+\gamma_5, \gamma^+$ and Eq. (19), we obtain the following expressions for the form factors $A_0(0)$ and $V(0)$ of the $Q' \to Q$ transition
\begin{align}
A_0(0) &= \int \frac{dx \, d^2k_{\perp} \phi_X^* \phi_V}{\sqrt{(A_V^2 + k_{\perp}^2)(A_X^2 + k_{\perp}^2)}} \left[ A_V A_X + (2x - 1) k_{\perp}^2 + \frac{2(m_{Q'} + m_Q)(1 - x)k_{\perp}^2}{W_V} \right], \quad (22) \\
V(0) &= \int \frac{dx \, d^2k_{\perp} \phi_X^* \phi_V}{\sqrt{(A_V^2 + k_{\perp}^2)(A_X^2 + k_{\perp}^2)}} (m_X + m_V)(1-x) \\
& \quad \left[ A_X + \frac{k_{\perp}^2}{W_V} + (1-x)(m_Q - m_{Q'})k_{\perp}^2 \right]. \quad (23)
\end{align}
where
\[ A_V = x m_q + (1 - x)m_{Q'} , \quad A_X = x m_q + (1 - x)m_Q \]
\[ W_V = M^V_0 + m_Q + m_{\bar{q}} , \quad \theta_V = \frac{d\phi_V}{dk^2} / \phi_V . \] (24)

Here, \( M^V_0 \) corresponds to Eq. (13) for the meson \( V \).

The meson wave functions \( \phi(x, k_{\perp}) \) are model dependent and difficult to obtain; often simple forms are assumed for them. One possibility is to use the wavefunction adopted in [11]
\[ \phi(x, k_{\perp}) = N \sqrt{x(1-x)} \exp \left( - \frac{M^2}{2w^2} \left[ x - \frac{1}{2} - \frac{m_Q^2 - m_{\bar{q}}^2}{2M^2} \right]^2 \right) \frac{\exp(-k^2_{\perp} / 2w^2)}{\sqrt{\pi w^2}} , \] (25)
where \( M \) is the mass of the meson. In [16] a Gaussian type of wavefunction was used,
\[ \phi(x, k_{\perp}) = N \sqrt{\frac{dk_z}{dx}} \exp \left( - \frac{k^2_{\perp}}{2\omega^2} \right) , \] (26)
where \( k_z \) is defined by \( x_1 M_0 = E_Q + k_z \) with \( E_Q = \sqrt{k^2_{\perp} + k_z^2 + m^2_Q} \). The parameter \( \omega \) in both Eqs. (25) and (26) should be of the order of \( \Lambda_{QCD} \). As we will see, the results are not too sensitive to the choice of wavefunction.

It is interesting to look at the general behavior of the wavefunction \( \phi(x, k_{\perp}) \) in the limit of \( m_Q \to \infty \). The distribution amplitude \( \int d^2k_{\perp} \phi(x, k_{\perp}) \) of a heavy meson, as is well known, should have a peak near \( x \simeq x_0 = \frac{m_Q}{m_V} \). As \( m_Q \) becomes larger the width of the peak decreases and \( x_0 \) comes closer to 1. The wavefunction \( \phi(x, k_{\perp}) \) vanishes if \( k^2_{\perp} \gg \Lambda^2_{QCD} \) and in the \( m_Q \to \infty \) limit peaks at \( x_0 \to 1 \) as the width of \( \phi(x, k_{\perp}) \to 0 \). Both wavefunctions listed have this feature.

3 The \( D^* \) decays

We now calculate the \( D^* \) decay rates in the light–front quark model. For \( D^* \to D\pi \), we first look at the form factor \( A_0(0) \) of the transition \( Q \to Q \). In this case Eq. (22)
reduces to

\[ A_0^{QQ}(0) = \int dx \, d^2k_\perp \phi_X^* \phi_V T_{A_0} \]

where

\[ T_{A_0} = \frac{A^2 + (2x - 1)k_\perp^2 + 4m_Q(1 - x)k_\perp^2}{A^2 + k_\perp^2} \]

and

\[ A = x \, m_{\bar{q}} + (1 - x) \, m_Q. \]

Generally \( T_{A_0} < 1 \) since it comes directly from the Melosh rotation. (For no Melosh rotation one would get \( T_{A_0} = 1 \).) How far \( T_{A_0} \) deviates from 1 depends on the relation among \( m_{\bar{Q}}, m_q, k_\perp^2 \), where the average value of \( k_\perp^2 \) is determined by the scale parameter in the wavefunction such as the \( \omega \) in Eqs. (25) and (26). Usually as \( < k_\perp^2 > \) becomes smaller \( T_{A_0} \) becomes closer to 1. If \( k_\perp = 0 \) in Eq. (28) then \( T_{A_0} = 1 \).

Let us consider two limiting cases for the transition quark \( Q \). If the transition quark \( Q \) is infinitely heavy while the spectator quark is light, \( x \to 1 \), and therefore the third term in the numerator of \( T_{A_0} \) vanishes. Thus

\[ T_{A_0} \to 1 \text{ and } A_0^{QQ}(0) \to \int dx \, d^2k_\perp \phi_X^* \phi_V \to 1, \]

as required by the heavy quark limit. Note in this case \( q^2 = q_{\text{max}}^2 = 0 \). If, as in the decays \( D^* \to D\pi \), the transition quark \( Q \) is light while the spectator quark is heavy, then in Eqs. (27) and (28), \( x \to 0 \) as \( m_c \to \infty \). In this case the third term in the numerator of \( T_{A_0} \) also vanishes, but now

\[ T_{A_0} \to \frac{A^2 - k_\perp^2}{A^2 + k_\perp^2}. \]

Since both \( k_\perp^2 \) and \( A^2 \) are of the same order, one can expect considerable deviation from 1 for \( A_0(0) \) in this case.
Obviously, the situation for \( g = A_0^{DD^*}(0) \) is quite close to the second case discussed above. Thus, one can clearly see why the light–front model gives a deviation of \( g \) from 1.

To calculate the radiative decays \( D^* \rightarrow D\gamma \), one can express \( \mu \) in Eq. (3) in terms of the form factor \( V(0) \)

\[
\mu = [e_c V^{cc}(0) + e_q V^{qq}(0)] \left( \frac{2}{m_D + m_{D^*}} \right)
\]

where \( q = u \) or \( d \) and \( e_c = \frac{2}{3}, e_u = \frac{2}{3}, e_d = -\frac{1}{3} \). \( V^{QQ}(0) \) is the \( V(0) \) form factor of the transition \( Q \rightarrow Q \) (\( Q = c, u \) or \( d \)). \( V^{QQ}(0) \) in the light–front model can be obtained from Eq. (23)

\[
V^{\bar{Q}Q}(0) = (m_D + m_{D^*}) \int \frac{dx \, d^2k_\perp \phi_{D^*}^\ast \phi_D}{A^2 + k_\perp^2} \left( 1 - x \right) \left\{ A + \frac{k_\perp^2}{W_{D^*}} \right\}.
\]

To calculate \( V^{cc}(0), V^{qq}(0), \) and \( g \), we use the two wavefunctions given in Eqs. (25) and (26) as the wavefunctions for the \( D \) and \( D^* \) mesons. We assume \( m_u = m_d = m \) and expect that the light–quark masses \( m \) in the wavefunctions to be around 0.25 \( \sim \) 0.30GeV. Similarly the scale parameter \( \omega \) in both wavefunctions should be around 0.40 \( \sim \) 0.50GeV. These values are similar to those taken in [11] and [16].

The dependence of \( V^{QQ}(0) \) on the quark mass \( m_Q \) is similar to that in the non–relativistic quark model [19, 20]: \( V^{QQ} = \frac{m_m + m_{D^*}}{2} \sqrt{\frac{m_{D^*} m_D}{m_Q}} \left( \frac{1}{m_Q} \right) \). That is \( V^{\bar{Q}Q}(0) \) decreases as \( m_Q \) increases and if \( m_Q \rightarrow \infty \), \( V^{\bar{Q}Q}(0) \rightarrow 1 \), as required by the heavy quark relation.

In the figure we show the dependence of the coupling \( g \) on the light quark mass \( m \) and \( \omega \) in the wavefunctions Eqs. (23) and (26). We choose \( m_c = 1.6\text{GeV} \); the coupling \( g \) has little dependence on \( m_c \). One can see that for a fixed \( \omega \), \( g \) increases as \( m \) increases. For \( m \) between 0.25 \( \sim \) 0.30GeV and \( \omega \) between 0.40 \( \sim \) 0.45GeV, \( g \) is between 0.55 \( \sim \) 0.65. To avoid adjusting parameters, we choose \( m = 0.25\text{GeV} \), and \( \omega = 0.4\text{GeV} \) for the calculation of decay rates. There is good agreement between the
calculated branching ratios and experiment. With this set of parameters $g = 0.6$ using Eq. (23) or Eq. (26). In the table, we show the corresponding results for $D^* \rightarrow D\pi$ and $D^* \rightarrow D\gamma$. We also show how the branching ratios change with a different choice of mass. Finally, we give the results of our calculation for the decay $D_{s}^{\ast +} \rightarrow D_{s}^{+}\gamma$ (for $m_{s} = 0.40\text{GeV}$ and $m_{s} = 0.50\text{GeV}$). This decay is usually assumed to be the dominant mode. In a recent paper [21] it has been argued that through isospin symmetry breaking there can also be a significant correlation between the branching ratios $\text{Br}(D_{s}^{\ast} \rightarrow D_{s}\pi^{0})$ and $\text{Br}(D_{s}^{\ast +} \rightarrow D_{s}^{+}\gamma)$.

4 Conclusion

In this letter we have presented a calculation of both strong and radiative $D^*$ decays using the relativistic light–front quark model. Our results are in good agreement with the 1992 CLEO II measurement. In particular the strong coupling $g$ is consistent with the experimental upper limit $\Gamma_{D^{\ast +}} < 131\text{keV}$. The key point is the relativistic treatment of the quark spin. The fact that the relativistic treatment is essential in explaining $g < 1$ is reminiscent of the similar situation in the nucleon axial–current form factor $g_{A}^{\text{nucleon}}$ where the non–relativistic quark model gives $g_{A}^{\text{nucleon}} = \frac{5}{3}$ and a relativistic treatment can get $g_{A}^{\text{nucleon}}$ naturally down to the measured value [22].

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Figure:

The dependence of the coupling $g$ on the light–quark mass $m$ and $\omega$ in Eqs. (25) and (26). The solid lines are the results of Eq. (26) and the dashed lines use Eq. (25). The upper solid and dashed lines correspond to $\omega = 0.40$ GeV, and the lower solid and dashed lines are for $\omega = 0.45$ GeV.
Table. Numerical results for $D^*$ decays. Also listed are the experimental results of CLEO II [3] and PDG [1]. All branching ratios are in %.

| Decay                  | Br$^{(a)}$ | Br$^{(b)}$ | Br$^{(a)}$ | Br$^{(b)}$ | CLEO II | PDG  |
|------------------------|------------|------------|------------|------------|---------|------|
| $D^{*+} \to D^0 \pi^+$ | 68.2       | 68.0       | 68.1       | 67.7       | 68.1±1.0±1.3 | 55±4 |
| $D^{*+} \to D^+ \pi^0$ | 31.6       | 31.5       | 31.6       | 31.4       | 30.8±0.4±0.8 | 27.2±2.5 |
| $D^{*+} \to D^+ \gamma$ | 0.3        | 0.5        | 0.4        | 0.9        | 1.1±1.4±1.6  | 18±4 |
| $D^{*0} \to D^0 \pi^0$ | 75.2       | 73.0       | 70.5       | 66.2       | 63.6±2.3±3.3 | 55±6 |
| $D^{*0} \to D^0 \gamma$ | 24.8       | 27.1       | 29.5       | 33.8       | 36.4±2.3±3.3 | 45±6 |

| Decay                  | $\Gamma$(keV)$^{(a)}$ | $\Gamma$(keV)$^{(a)}$ | $\Gamma$(keV)$^{(b)}$ | $\Gamma$(keV)$^{(b)}$ |
|------------------------|------------------------|------------------------|------------------------|------------------------|
| $D^{*+} \to$ total     | 122.8                  | 123.0                  | 104.9                  | 102.1                  |
| $D^{*0} \to$ total     | 74.0                   | 76.3                   | 67.4                   | 69.5                   |
| $D^{*+}_s \to D^*_s \gamma$ | 0.1                    | 0.1                    | 0.2                    | 0.3                    |

$a): \omega = 0.4$GeV, $m = m_u = m_d = 0.30$GeV, $m_s = 0.50$GeV.

$b): \omega = 0.4$GeV, $m = m_u = m_d = 0.25$GeV, $m_s = 0.40$GeV.
This figure "fig1-1.png" is available in "png" format from:

http://arxiv.org/ps/hep-ph/9406300v1