

Couplings between pion and charmed mesons

Hungchong Kim^a, Su Hounng Lee

Institute of Physics and Applied Physics, Yonsei University, Seoul 120-749, Korea

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Abstract. We compute the couplings $DD^*\pi$, $D_1D^*\pi$, $D^*D^*\pi$ and $D_1D_1\pi$ using QCD sum rules. These couplings are important inputs in the meson-exchange model calculations used to estimate the amount of J/ψ absorption due to pions and rho mesons in heavy-ion collisions. Our sum rules are constructed at the first order in the pion momentum p_μ , which give the couplings that are not trivially related to the soft-pion theorem. Our calculated couplings, which somewhat depend upon the values of the heavy meson decay constants, are $g_{DD^*\pi} = 8.2 \pm 0.1$, $g_{D_1D^*\pi} = 15.8 \pm 2$ GeV, $g_{D^*D^*\pi} = 0.3 \pm 0.03$ and $g_{D_1D_1\pi} = 0.17 \pm 0.04$.

1 Introduction

J/ψ suppression [1] seems to be one of the most promising signals for QGP formation in RHIC. Indeed the recent data at CERN [2] show an anomalous suppression of J/ψ formation, which seems to be a consequence of QGP formation [3]. However, before coming to a conclusion, one has to estimate the amount of J/ψ suppression due to hadronic final state interactions.

Consequently, there have been a number of works, using various models [4–13], calculating the J/ψ absorption cross section by light mesons. However, the estimate varies by an order of magnitude near threshold. At this stage, it is necessary to probe all model calculations further to spell out their corrections and uncertainties.

In the effective meson-exchange model approaches [9–13], important ingredients are the couplings between open charm mesons and light mesons. The precise determination of the couplings reduces the uncertainties in the calculation of the dissociation process. Moreover, a complete set of low-lying open charm mesons, $D(1870)$, $D^*(2010)$ and $D_1(2420)$, has to be included. This is especially necessary in order to probe the cross section above the threshold. To provide the basic building blocks for such a model calculation, we use QCD sum rules [14,15] and calculate the couplings $DD^*\pi$, $D^*D^*\pi$, $D_1D_1\pi$ and $D_1D^*\pi$. These can be used to improve the existing effective model calculations, which can be applied to a future calculation [16] of the J/ψ dissociation process.

At present, there are two approaches in the literature to calculate the coupling in the QCD sum rule approach: light-cone QCD sum rules (LCQSR) [17,19,18] relying on the operator product expansion near the light cone, and the conventional QCD sum rules [20–22] based on the short-distance expansion. Predictions from LCQSR heavily depend on the twist-2 pion wave function at the middle

point, whose value, however, has been at the core of debates for a long time [23]. Furthermore, the duality issue in constructing the continuum needs to be carefully considered [24]. Instead, the conventional QCD sum rules [20–22] do not suffer from such an uncertainty as the QCD parameters appearing in this approach are determined from the low-energy theorems such as PCAC and the soft-pion theorem. Though QCD duality again needs to be applied carefully in these sum rules [25,26], the uncertainties from the QCD side can be substantially reduced.

In this work, we provide a systematic approach to calculate the couplings using the conventional QCD sum rules. Our sum rules will be constructed at the first order in the pion momentum $\mathcal{O}(p_\mu)$. We improve the previous sum rule calculations of the $DD^*\pi$ and $D^*D^*\pi$ couplings [20–22]. In particular, QCD duality will be correctly applied according to the recent suggestion in [25,26]. Also in constructing a sum rule for $DD^*\pi$, we advocate the use of a different structure function. We then construct similar sum rules for the couplings, $D_1D_1\pi$ and $D_1D^*\pi$, which may be important for future calculation of the J/ψ dissociation process accompanying the D_1 meson [16].

2 OPE for general correlation function

In this section, we schematically describe a general procedure to perform the operator product expansion (OPE) of the general correlation function with a pion,

$$\begin{aligned} \Pi(p, q) = & i \int d^4x e^{iq \cdot x} \\ & \times \langle 0 | T \{ \bar{d}(x) \Gamma_1 c(x), \bar{c}(0) \Gamma_2 u(0) \} | \pi(p) \rangle. \end{aligned} \quad (1)$$

Γ_1 and Γ_2 denote specific gamma matrices corresponding to the coupling of concern. In later sections, we will use this general prescription, by simple replacements, to

^a e-mail: hung@phya.yonsei.ac.kr

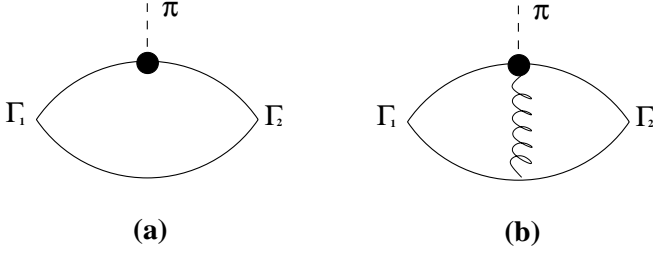


Fig. 1a,b. The OPE diagrams considered in this work. The blob in **a** denotes the quark–antiquark component with a pion and the blob in **b** denotes the quark–antiquark–gluon component with a pion

calculate the OPE for the correlation function of concern. For instance, to calculate the $DD^*\pi$ coupling, we will have $\Gamma_1 = \gamma_\mu$ and $\Gamma_2 = i\gamma_5$.

To calculate the coupling that is not trivially related by chiral symmetry, we will consider the correlation function at the first order of the pion momentum p_μ in its expansion. The soft-pion limit of the correlation function is just a chiral rotation of a vacuum correlation function (i.e. *without the pion*), which provides a coupling that is trivially related to the vacuum correlation function. For example, in the case of the pion–nucleon coupling [27], the soft-pion limit leads to a Goldberger–Treiman relation with $g_A = 1$, which gives a coupling 30% lower than its experimental value. To determine the coupling more precisely, one needs to go beyond the soft-pion limit.

First, we restrict ourselves to the OPE diagram shown in Fig. 1a. For this diagram, we can rewrite the correlation function in the form

$$-i \int \frac{d^4k}{(2\pi)^4} \text{Tr} \left[i \frac{\not{q} - \not{k} + m_c}{(q-k)^2 - m_c^2} \Gamma_2 D_{aa}(k, p) \Gamma_1 \right]. \quad (2)$$

Roman indices denote colors. The free c -quark propagator has been used in obtaining this. $D_{ab}(k, p)$, which is shown by the blob in the figure, denotes the quark–antiquark component with a pion. This can be separated into three pieces depending on the Dirac matrices involved,

$$D_{ab}(k, p) = \delta_{ab} \left[i\gamma_5 A(k, p) + \gamma_\alpha \gamma_5 B^\alpha(k, p) + \gamma_5 \sigma_{\alpha\beta} C^{\alpha\beta}(k, p) \right]. \quad (3)$$

The three invariant functions of k, p are defined by

$$\begin{aligned} A(k, p) &= \frac{1}{12} \sum_{n=0}^{\infty} \frac{1}{n!} \langle 0 | \bar{d} \overleftarrow{D}_{\alpha_1} \cdots \overleftarrow{D}_{\alpha_n} i\gamma_5 u | \pi(p) \rangle \\ &\quad \times (2\pi)^4 \frac{\partial^n}{i\partial k_{\alpha_1} \cdots i\partial k_{\alpha_n}} \delta^{(4)}(k), \\ B^\alpha(k, p) &= \frac{1}{12} \sum_{n=0}^{\infty} \frac{1}{n!} \langle 0 | \bar{d} \overleftarrow{D}_{\alpha_1} \cdots \overleftarrow{D}_{\alpha_n} \gamma^\alpha \gamma_5 u | \pi(p) \rangle \\ &\quad \times (2\pi)^4 \frac{\partial^n}{i\partial k_{\alpha_1} \cdots i\partial k_{\alpha_n}} \delta^{(4)}(k), \\ C^{\alpha\beta}(k, p) &= -\frac{1}{24} \sum_{n=0}^{\infty} \frac{1}{n!} \langle 0 | \bar{d} \overleftarrow{D}_{\alpha_1} \cdots \overleftarrow{D}_{\alpha_n} \gamma_5 \sigma^{\alpha\beta} u | \pi(p) \rangle \end{aligned}$$

$$\times (2\pi)^4 \frac{\partial^n}{i\partial k_{\alpha_1} \cdots i\partial k_{\alpha_n}} \delta^{(4)}(k). \quad (4)$$

Since we are constructing the sum rules at the order $\mathcal{O}(p_\mu)$, we need to evaluate the pion matrix elements at $\mathcal{O}(p_\mu)$. Noting that the pion matrix elements involved are symmetric under exchanges of any pair of indices $\alpha_1 \cdots \alpha_n$ and using the soft-pion theorem and PCAC, we straightforwardly obtain

$$\begin{aligned} \langle 0 | \bar{d} \overleftarrow{D}_{\alpha_1} i\gamma_5 u | \pi(p) \rangle &= \frac{\langle \bar{q}q \rangle}{f_\pi} i p_{\alpha_1}, \\ \langle 0 | \bar{d} \overleftarrow{D}_{\alpha_1} \overleftarrow{D}_{\alpha_2} \overleftarrow{D}_{\alpha_3} i\gamma_5 u | \pi(p) \rangle &= \frac{im_0^2 \langle \bar{q}q \rangle}{12f_\pi} (p_{\alpha_1} g_{\alpha_2 \alpha_3} + p_{\alpha_2} g_{\alpha_1 \alpha_3} + p_{\alpha_3} g_{\alpha_1 \alpha_2}), \\ \langle 0 | \bar{d} \gamma^\alpha \gamma_5 u | \pi(p) \rangle &= i p^\alpha f_\pi, \\ \langle 0 | \bar{d} \overleftarrow{D}_{\alpha_1} \overleftarrow{D}_{\alpha_2} \gamma_\alpha \gamma_5 u | \pi(p) \rangle &= i f_\pi \delta^2 \left[p_\alpha g_{\alpha_1 \alpha_2} \frac{5}{18} - (p_{\alpha_2} g_{\alpha \alpha_1} + p_{\alpha_1} g_{\alpha \alpha_2}) \frac{1}{18} \right], \\ \langle 0 | \bar{d} \overleftarrow{D}_{\alpha_1} \gamma_5 \sigma_{\alpha\beta} u | \pi(p) \rangle &= i (p_\alpha g_{\beta \alpha_1} - p_\beta g_{\alpha \alpha_1}) \frac{\langle \bar{q}q \rangle}{3f_\pi}, \\ \langle 0 | \bar{d} \overleftarrow{D}_{\alpha_1} \overleftarrow{D}_{\alpha_2} \overleftarrow{D}_{\alpha_3} \gamma_5 \sigma_{\alpha\beta} u | \pi(p) \rangle &= \frac{im_0^2 \langle \bar{q}q \rangle}{36f_\pi} [p_\alpha (g_{\alpha_1 \alpha_2} g_{\alpha_3 \beta} + g_{\alpha_1 \alpha_3} g_{\alpha_2 \beta} + g_{\alpha_3 \alpha_2} g_{\alpha_1 \beta}) \\ &\quad - (\alpha \leftrightarrow \beta)]. \end{aligned} \quad (5)$$

Here m_0^2 and δ^2 are defined via

$$\begin{aligned} \langle \bar{q} D^2 q \rangle &= \frac{m_0^2}{2} \langle \bar{q}q \rangle, \\ \langle 0 | \bar{d} g_s \tilde{G}^{\alpha\beta} \gamma_\beta u | \pi(p) \rangle &= i \delta^2 f_\pi p^\alpha. \end{aligned} \quad (6)$$

Up to twist-5, these are all the possibilities coming from the expansion of the quark–antiquark components at the order $\mathcal{O}(p_\mu)$.

The additional contribution to the OPE is shown by Fig. 1b where one gluon emitted from the c -quark propagator is combined with the quark–antiquark component. Specifically, the c -quark propagator with one gluon being attached is given by [28]

$$-\frac{g_s \mathcal{G}_{\alpha\beta}}{2(k^2 - m_c^2)^2} [k^\alpha \gamma^\beta - k^\beta \gamma^\alpha + (\not{k} + m_c) i \sigma^{\alpha\beta}], \quad (7)$$

where $\mathcal{G}_{\alpha\beta} = t^A G_{\alpha\beta}$. The color matrices t^A are normalized via $\text{Tr}(t^A t^B) = \delta^{AB}/2$. Taking the gluon stress tensor into the quark–antiquark component, one can write down the correlation function into the form

$$\begin{aligned} \Pi(p, q) &= 2i \int \frac{d^4k}{(2\pi)^4} \\ &\quad \times \text{Tr} \left\{ \frac{2(q-k)_\theta \gamma_\delta + (\not{q} - \not{k} + im_c) i \sigma_{\theta\delta}}{[(q-k)^2 - m_c^2]^2} \right. \\ &\quad \left. \times \Gamma_2 [\gamma_5 \sigma_{\rho\lambda} B^{\rho\lambda\theta\delta}(k, p) + \gamma^\tau \epsilon^{\theta\delta\alpha\beta} D_{\tau\alpha\beta}(k, p)] \Gamma_1 \right\}. \end{aligned} \quad (8)$$

At the order $\mathcal{O}(p_\mu)$, the two functions appearing above are given by

$$B^{\rho\lambda\theta\delta} = -\frac{m_0^2 \langle \bar{q}q \rangle}{12 \times 32 f_\pi} (g^{\rho\theta} g^{\lambda\delta} - g^{\rho\delta} g^{\lambda\theta}) p^\alpha (2\pi)^4 \times \frac{\partial}{\partial k_\alpha} \delta^{(4)}(k), \quad (9)$$

$$D_{\tau\alpha\beta} = -\frac{i\delta^2 f_\pi}{3 \times 32} (p_\alpha g_{\tau\beta} - p_\beta g_{\tau\alpha}) (2\pi)^4 \delta^{(4)}(k). \quad (10)$$

Another function involving the pion matrix element of the form

$$\langle 0 | \bar{d} \gamma_5 \gamma_\tau g_s \mathcal{G}_{\theta\delta} u | \pi(p) \rangle \quad (11)$$

is of the order $\mathcal{O}(m_q)$, which therefore has been neglected.

Once the Dirac matrices $\Gamma_i (i = 1, 2)$ are given, we can straightforwardly calculate the corresponding OPE from (2) and (8). These give the OPE up to twist-5 at the order $\mathcal{O}(p_\mu)$. Below, we will use this formalism to calculate the couplings $DD^*\pi$, $D^*D^*\pi$, $D_1D_1\pi$, and $D^*D_1\pi$ by choosing appropriate Dirac matrices for them.

3 Sum rule for $DD^*\pi$

In this section, we construct a sum rule for the $DD^*\pi$ coupling using the correlation function

$$\Pi_\mu(p, q) = i \int d^4x e^{iq \cdot x} \langle 0 | T \{ \bar{d}(x) \gamma_\mu c(x), \bar{c}(0) i \gamma_5 u(0) \} | \pi(p) \rangle. \quad (12)$$

As two momenta are involved in the correlation function, one can separate the correlation function in the following two pieces:

$$\Pi_\mu = F_1(p, q) p_\mu + F_2(q, p) q_\mu. \quad (13)$$

We construct a sum rule for $F_1(p = 0, q)$, which gives a coupling determined at the order $\mathcal{O}(p_\mu)$. Defining the $D^*D\pi$ coupling by [17]

$$\langle D^*(q) | D(q-p) \pi(p) \rangle = g_{D^*D\pi} p \cdot \epsilon, \quad (14)$$

and using

$$\begin{aligned} \langle D | \bar{c} i \gamma_5 u | 0 \rangle &= \frac{m_D^2 f_D}{m_c}, \\ \langle 0 | \bar{d} \gamma_\mu c | D^* \rangle &= m_{D^*} f_{D^*} \epsilon_\mu, \end{aligned} \quad (15)$$

the low-lying pole contribution to $F_1(p = 0, q)$ is given by

$$-\frac{m_D^2 m_{D^*} f_D f_{D^*} g_{DD^*\pi}}{m_c (q^2 - m_{D^*}^2) (q^2 - m_D^2)}. \quad (16)$$

A slightly different sum rule using the same correlation function can be found in [20, 21]. Specifically, the correlation function in that approach was decomposed into

$$\Pi_\mu = A p_\mu + B(2q - p)_\mu \quad (17)$$

and a sum rule was constructed for the function A . Note that, by comparing this with (13), one can immediately see that $A = F_1 + F_2/2$. Thus, in the expansion in terms of the external momentum p_μ , the function A involves the term at the zeroth order in p_μ (i.e. F_2), which can be trivially obtained from a vacuum correlation function via the soft-pion theorem, as well as the term of the first order in p_μ (i.e., F_1). To avoid the trivial contribution obtained from the soft-pion theorem, we choose to work with the decomposition of (13). Furthermore, as mentioned in [17], the function A can contain some contribution from a possible resonance (scalar particle D_0), which can give an additional uncertainty in the prediction.

Following the general strategy given in Sect. 2, we obtain the OPE for the correlation function,

$$\begin{aligned} F_1(q, p = 0) &= \frac{1}{q^2 - m_c^2} \left[m_c f_\pi - \frac{2 \langle \bar{q}q \rangle}{3 f_\pi} \left(2 - \frac{m_c^2}{q^2 - m_c^2} \right) \right. \\ &\quad \left. - \frac{10 \delta^2 f_\pi m_c}{9 q^2 - m_c^2} \left(1 + \frac{m_c^2}{q^2 - m_c^2} \right) \right] \\ &\quad + \frac{m_0^2 \langle \bar{q}q \rangle}{6 f_\pi} \left[\frac{5}{6(q^2 - m_c^2)^2} + \frac{4m_c^2}{3(q^2 - m_c^2)^3} \right. \\ &\quad \left. - \frac{4m_c^4}{(q^2 - m_c^2)^4} \right]. \end{aligned} \quad (18)$$

Note that the leading OPE has a simple pole in q^2 . According to QCD duality, higher resonance contributions lying along the positive q^2 are matched with the imaginary part of the OPE above a certain threshold S_0 which is taken much higher than the low-lying pole. Since the simple pole structure $1/(q^2 - m_c^2)$ does not have non-analytic structures in the duality region ($q^2 \geq S_0$), it should not pick up the continuum contribution. However, it is a common practice that the double-variable dispersion relation is used to obtain a spectral density for a given OPE. Then by naively restricting the dispersion integral below the continuum threshold, one picks up the continuum contribution even from the OPE of the form $1/(q^2 - m_c^2)$. In this prescription, the continuum contribution is a simple (and unphysical) pole at the continuum threshold [25, 26, 24]. (See for example (3.12) of [22].) This pole at the continuum threshold does not resemble at all the higher resonance contributions lying along the positive q^2 . In fact, this pole at the continuum is mathematically spurious. To illustrate this in detail, let us determine the spectral density for the OPE $1/(q^2 - m_c^2)$ from the double dispersion relation when the external momentum is zero,

$$\frac{1}{q^2 - m_c^2} = \int_0^\infty ds \frac{b(s)}{(s - q^2)^2}. \quad (19)$$

Under the successive Borel transformations [26], one can determine the spectral function

$$b(s) = -\theta(s - m_c^2). \quad (20)$$

When we put it back to the double dispersion relation, we have to reproduce the OPE $1/(q^2 - m_c^2)$. Anything

additional to it is mathematically spurious. Using (20) in (19) and doing the integration by part, we obtain

$$\int_0^\infty ds \frac{-\theta(s - m_c^2)}{(s - q^2)^2} = \frac{1}{q^2 - m_c^2} + \frac{\theta(s - m_c^2)}{s - q^2} \Big|_0^\infty. \quad (21)$$

The second term is mathematically spurious as it is additional to the one that we had started with. But when it is restricted by the continuum threshold S_0 , the upper limit is changed to S_0 and the second term has a pole at the continuum threshold. However, since the pole comes from the spurious term, its contribution to the sum rule is spurious.

Nonetheless, one may argue from intuition that the continuum contribution should be present as the current can couple to higher resonances. In fact, it may be possible to build such a contribution if one uses a more sophisticated current than the simple current of the form $\bar{q}\Gamma c$. We believe that the absence of the continuum is due to a limitation of the current of the form $\bar{q}\Gamma c$. But we believe that, as we demonstrated above, it is ad hoc to build the continuum contribution from the OPE of the form $1/(q^2 - m_c^2)$.

We now match (16) with (18) to get a sum rule for the $D^*D\pi$ coupling. For simplicity, we neglect the mass difference between D^* and D and set them to $m_{D^*} = m_D = m_{av} \equiv (m_{D^*} + m_D)/2$. Under the Borel transformation (with the Borel mass M^2), the final sum rule reads

$$\begin{aligned} g_{DD^*\pi} f_D f_{D^*} + T_1 M^2 & \quad (22) \\ = \frac{m_c}{m_{av}^3} M^2 e^{(m_{av}^2 - m_c^2)/M^2} & \left\{ m_c f_\pi - \frac{4}{3f_\pi} \langle \bar{q}q \rangle \right. \\ + \left[-\frac{2}{3f_\pi} m_c^2 \langle \bar{q}q \rangle + \frac{10}{9} \delta^2 f_\pi m_c - \frac{5}{36f_\pi} m_0^2 \langle \bar{q}q \rangle \right] & \frac{1}{M^2} \\ + \left[-\frac{5}{9} \delta^2 f_\pi m_c^3 + \frac{1}{9f_\pi} m_0^2 m_c^2 \langle \bar{q}q \rangle \right] & \frac{1}{M^4} + \frac{m_0^2 m_c^4 \langle \bar{q}q \rangle}{9f_\pi M^6} \left. \right\}. \end{aligned}$$

Here T_1 denotes the transitions from the low-lying resonance to higher resonances. We will linearly fit the RHS within a Borel window to determine the coupling as well as the transition strength T_1 .

4 Sum rule for $D^*D^*\pi$

We now construct a sum rule for the $D^*D^*\pi$ coupling. The $D^*D^*\pi$ sum rule can be constructed from the correlation function (by setting $\Gamma_1 = \gamma_\mu$ and $\Gamma_2 = \gamma_\nu$ in (1)). We have

$$\begin{aligned} H_{\mu\nu}(p, q) & = i \int d^4x e^{iq \cdot x} \quad (23) \\ & \times \langle 0 | T \{ \bar{d}(x) \gamma_\mu c(x), \bar{c}(0) \gamma_\nu u(0) \} | \pi(p) \rangle. \end{aligned}$$

Saturating the correlation function by the D^* intermediate state and introducing the coupling via

$$\begin{aligned} \langle D^*(q, \epsilon_2) | \pi(p) D^*(q - p, \epsilon_1) \rangle & \\ = i \frac{2}{f_\pi} g_{D^*D^*\pi} \epsilon_{\alpha\beta\mu\nu} \epsilon_1^\alpha \epsilon_2^\beta p^\mu q^\nu, & \quad (24) \end{aligned}$$

the low-lying pole contribution at the first order in p_μ is given by [22]

$$-\frac{2g_{D^*D^*\pi} f_{D^*}^2 m_{D^*}^2}{f_\pi (q^2 - m_{D^*}^2)^2} \epsilon_{\alpha\beta\mu\nu} p^\mu q^\nu. \quad (25)$$

The OPE part can be computed by following the general prescription described in Sect. 2. After taking out the common factor of $\epsilon_{\alpha\beta\mu\nu} p^\mu q^\nu$, we obtain the OPE side

$$\begin{aligned} \frac{f_\pi}{q^2 - m_c^2} + \frac{2}{3f_\pi} \frac{m_c \langle \bar{q}q \rangle}{(q^2 - m_c^2)^2} + \frac{8}{9} \frac{f_\pi \delta^2}{(q^2 - m_c^2)^2} & \\ - \frac{10}{9} \frac{f_\pi \delta^2 m_c^2}{(q^2 - m_c^2)^3} - \frac{2}{3f_\pi} \frac{m_c^3 m_0^2 \langle \bar{q}q \rangle}{(q^2 - m_c^2)^4}. & \quad (26) \end{aligned}$$

The terms involving m_0^2 are different from the one in [22]. By matching the two sides, we obtain the sum rule for the $D^*D^*\pi$ coupling

$$\begin{aligned} g_{D^*D^*\pi} f_{D^*}^2 + T_2 M^2 & \\ = \frac{f_\pi}{2m_{D^*}^2} e^{(m_{D^*}^2 - m_c^2)/M^2} \left[f_\pi M^2 - \frac{2}{3f_\pi} m_c \langle \bar{q}q \rangle - \frac{8}{9} f_\pi \delta^2 \right. & \\ \left. - \frac{5}{9} f_\pi \delta^2 \frac{m_c^2}{M^2} + \frac{m_c^3 m_0^2 \langle \bar{q}q \rangle}{9f_\pi M^4} \right]. & \quad (27) \end{aligned}$$

Here, T_2 denotes the transitions from $D^* \rightarrow$ higher resonance states.

5 Sum rule for $D_1D_1\pi$

For the coupling $D_1D_1\pi$, we use the correlation function involving axial-vector currents,

$$i \int d^4x e^{iq \cdot x} \langle 0 | T \{ \bar{d}(x) \gamma_\mu \gamma_5 c(x), \bar{c}(0) \gamma_\nu \gamma_5 u(0) \} | \pi(p) \rangle. \quad (28)$$

Comparing this correlation function with (23), one can easily see that the OPE for this correlation function can be obtained by replacing $m_c \rightarrow -m_c$ in (26). We introduce the $D_1D_1\pi$ coupling similarly as (24). Then it is a simple matter to construct a sum rule for the $D_1D_1\pi$ coupling. Namely, by replacing $m_c \rightarrow -m_c$ in (27), we have

$$\begin{aligned} g_{D_1D_1\pi} f_{D_1}^2 + T_3 M^2 & \\ = \frac{f_\pi}{2m_{D_1}^2} e^{(m_{D_1}^2 - m_c^2)/M^2} \left[f_\pi M^2 + \frac{2}{3f_\pi} m_c \langle \bar{q}q \rangle - \frac{8}{9} f_\pi \delta^2 \right. & \\ \left. - \frac{5}{9} f_\pi \delta^2 \frac{m_c^2}{M^2} - \frac{m_c^3 m_0^2 \langle \bar{q}q \rangle}{9f_\pi M^4} \right]. & \quad (29) \end{aligned}$$

Again, T_3 denotes the transitions from $D_1 \rightarrow$ higher resonance states.

6 Sum rule for $D_1D^*\pi$

The $D_1D^*\pi$ coupling can be calculated from the correlation function

$$i \int d^4x e^{iq \cdot x} \langle 0 | T \{ \bar{d}(x) \gamma_\mu c(x), \bar{c}(0) \gamma_\nu \gamma_5 u(0) \} | \pi(p) \rangle. \quad (30)$$

[19, 18] calculated the coupling using the light-cone QCD sum rules. Both references considered the structure function proportional to $g_{\mu\nu}$, but the OPE in either approach is different. Here we choose to work with the structure function proportional to $q_\mu p_\nu + q_\nu p_\mu$, which turns out to be rather simple.

In constructing the phenomenological side, we follow [19] where the structure proportional to $q_\mu p_\nu + q_\nu p_\mu$ is given by

$$g_{D_1 D^* \pi} \frac{m_{D_1}}{m_{D^*}} \frac{f_{D^*}}{q^2 - m_{D^*}^2} \frac{f_{D_1}}{q^2 - m_{D_1}^2}. \quad (31)$$

Note that $g_{D_1 D^* \pi}$ has one dimension due to the way that the coupling is introduced in [19], which is in contrast to the other dimensionless couplings. The OPE side can be calculated straightforwardly. It takes the simple form of

$$-f_\pi \left[\frac{1}{q^2 - m_c^2} - \frac{10}{9} \frac{m_c^2}{(q^2 - m_c^2)^3} \right]. \quad (32)$$

The term containing δ^2 from the expansion of $A(k, p)$ is canceled with the similar term coming from $D_{\tau\alpha\beta}$. Setting $m_{D^*} = m_{D_1} = \bar{m} \equiv (m_{D^*} + m_{D_1})/2$ and matching the two sides, we obtain

$$g_{D_1 D^* \pi} f_{D^*} f_{D_1} + T_4 M^2 = f_\pi e^{(\bar{m}^2 - m_c^2)/M^2} \left[M^2 - \frac{5}{9} \frac{m_c^2}{M^2} \right]. \quad (33)$$

7 Analysis and results

In our analysis, we use the following set of QCD parameters:

$$\begin{aligned} m_0^2 &= 0.8 \text{ GeV}^2, & \langle \bar{q}q \rangle &= (-0.24 \text{ GeV})^3, \\ \delta^2 &= 0.2 \text{ GeV}^2, & m_c &= 1.34 \text{ GeV}, & f_\pi &= 131 \text{ MeV}. \end{aligned} \quad (34)$$

For the hadron masses, we use [29]

$$m_D = 1.87 \text{ GeV}, \quad m_{D^*} = 2.01 \text{ GeV}, \quad m_{D_1} = 2.42 \text{ GeV}. \quad (35)$$

We plot the Borel curves for the couplings $DD^*\pi$ and $D_1 D^* \pi$ in Figs. 2 and 3 using (22) and (33) respectively. The corresponding curves for the $D^* D^* \pi$ and $D_1 D_1 \pi$ couplings are shown in Fig. 3. To get the couplings, we need to fit each curve with a straight line within an appropriately chosen Borel window. The intersection between the best fitting curve and the vertical axis at $M^2 = 0$ gives the values $f_D f_{D^*} g_{DD^* \pi}$, $f_{D_1} f_{D^*} g_{D_1 D^* \pi}$, $f_{D^*}^2 g_{D^* D^* \pi}$ and $f_{D_1}^2 g_{D_1 D_1 \pi}$. All the curves are well fitted with a straight line above minimum Borel mass, which depending on the sum rules, ranges from 2–3 GeV^2 . In that region, the higher dimensional terms are suppressed. When we shift the Borel window by 0.5 GeV^2 , the results reduce by 10%.

Table 1 shows the best fit values and the chosen Borel window. For $f_D f_{D^*} g_{DD^* \pi}$, our value is substantially smaller than 0.51 GeV^2 obtained from the light-cone QCD

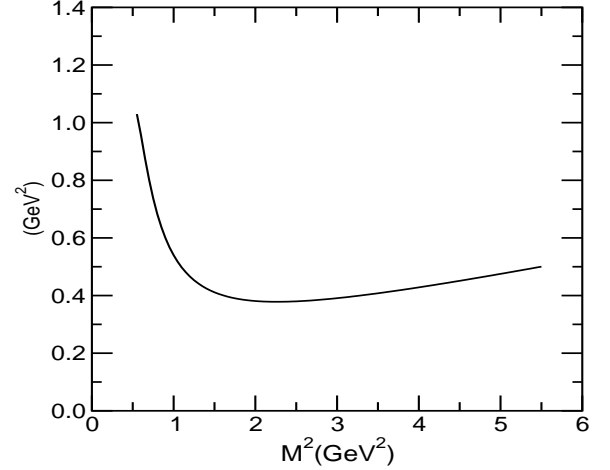


Fig. 2. The Borel curves for the $DD^*\pi$ coupling from (22). Here the RHS of the equation is plotted with respect to the Borel mass M^2

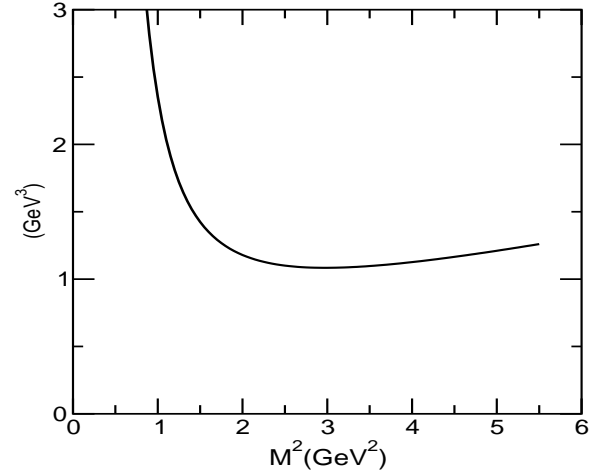


Fig. 3. The Borel curves for the $D_1 D^* \pi$ coupling from (33). Here the RHS of the equation is plotted with respect to the Borel mass M^2

Table 1. The best fitted values for the couplings are listed here along with the chosen Borel window. An overall shift of the Borel window by 0.5 GeV^2 to the higher mass region gives 10% error

Borel window (GeV^2)	fitted value
$2.5 \leq M^2 \leq 3.5$	$f_D f_{D^*} g_{DD^* \pi} = 0.31 \text{ GeV}^2$
$3 \leq M^2 \leq 4$	$f_{D_1} f_{D^*} g_{D_1 D^* \pi} = 0.948 \text{ GeV}^3$
$2 \leq M^2 \leq 3$	$f_{D^*}^2 g_{D^* D^* \pi} = 0.016 \text{ GeV}^2$
$2 \leq M^2 \leq 3$	$f_{D_1}^2 g_{D_1 D_1 \pi} = 0.012 \text{ GeV}^2$

sum rule analysis [17]. The origin of the difference may be traced back to the use of the asymptotic pion wave functions. In the light-cone sum rules, the values of the twist-2 and twist-3 wave functions at the middle point, $\varphi_\pi(1/2) \sim 1.2$ and $\varphi_p(1/2) \sim 1.5$, enter as leading terms in the OPE. But in our sum rules, these middle points are replaced by the integrated strengths of the wave func-

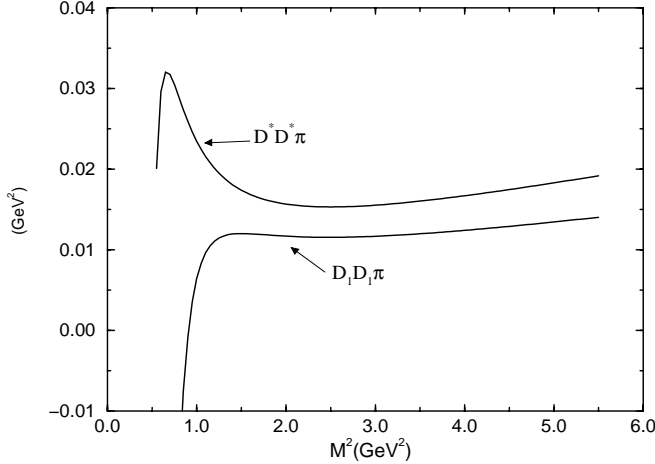


Fig. 4. The Borel curves for the $D^*D^*\pi$ and $D_1D_1\pi$ couplings given in (27) and (29). Here the RHS of the equations are plotted with respect to the Borel mass M^2

tions $\int_0^1 du \varphi_\pi(u) = 1$, $\int_0^1 du \varphi_p(u) = 1$, which are well fixed by low-energy theorems. Though our approach is different from the similar calculations in [20,22], our result for $f_D f_{D^*} g_{DD^*\pi}$ agrees with them. Also our value of $f_{D^*}^2 g_{D^*D^*\pi}$ agrees with the results in [18]. Our result for $f_{D_1} f_{D^*} g_{D_1D^*\pi}$, 0.948 GeV^3 , is somewhat larger than 0.68 GeV^3 from [19], which is obtained from the structure function being proportional to $g_{\mu\nu}$.

To get the couplings, we need to determine f_D , f_{D_1} and f_{D^*} . One may calculate these using QCD sum rules of the two-point correlation function in the vacuum. Currently these values are not known precisely. According to [18], they are $f_D = 170 \text{ MeV}$, $f_{D^*} = 220 \text{ MeV}$ and $f_{D_1} = 240 \text{ MeV}$. A somewhat different set of the decay constants can be found in [19], $f_D = 160 \text{ MeV}$, $f_{D^*} = 240 \text{ MeV}$ and $f_{D_1} = 300 \text{ MeV}$. Using these, we obtain the coupling constants

$$\begin{aligned} g_{DD^*\pi} &= 8.29(8.07), & g_{D_1D^*\pi} &= 17.95(13.6) \text{ GeV}, \\ g_{D^*D^*\pi} &= 0.33(0.278), & g_{D_1D_1\pi} &= 0.208(0.133). \end{aligned} \quad (36)$$

where the numbers (the ones in parentheses) are obtained by using the decay constants given in [18] ([19]). To make a more precise prediction, one may need to determine the decay constants precisely. To summarize, we present in Table 2 our results in comparison with the other previous calculations.

These couplings are important ingredients for estimating the absorption cross section of J/ψ by π mesons. Up to now meson-exchange models in the calculation of the absorption cross section are based on an effective chiral lagrangian with only pseudoscalars (D, π) and vector mesons ($J/\psi, D^*$). We believe that the addition of the axial partner of D^* , namely the D_1 meson, will reduce the value for the existing calculation. The reason is this. Consider the dissociation of J/ψ by pions into D and D^* . In the existing meson-exchange calculations, the diagrams contributing to the process are the contributions from the s -channel D meson, the t -channel D^* mesons and the direct four-point coupling, which gives a non-trivial contri-

Table 2. Comparison of our results with the other published results. The results of [17–19] are from light-cone QCD sum rules, the results from [20,22] are from the conventional sum rules, and the result of [30] is from the three-point sum rules

	$g_{DD^*\pi}$	$g_{D_1D^*\pi}$ (GeV)	$g_{D^*D^*\pi}$	$g_{D_1D_1\pi}$
this work	8.2 ± 0.1	15.8 ± 2	0.3 ± 0.03	0.17 ± 0.04
[17]	12.5 ± 1			
[18]	11.85 ± 2.1		0.31 ± 0.08	
[19]		10 ± 2		
[20,22]	9 ± 2		0.35 ± 0.08	
[30]	5.7 ± 4			

bution. The form of the couplings are obtained from the chiral SU(4) lagrangian with vector mesons introduced in the massive Yang–Mills approach. We believe that the addition of the D_1 meson will partly cancel the contribution from the direct four-point coupling. This is so because this is precisely how the Adler consistency condition for the ρ - π forward scattering amplitude is obtained in the massive Yang Mills approach [31]. Namely, one has to introduce the axial partner of ρ , the a_1 meson, whose contribution in the s -channel will cancel the direct four-point coupling of the $\rho\rho\pi\pi$ and make the amplitude vanish in the soft-pion limit. A similar cancellation will occur between the direct four-point coupling of J/ψ - π - D - D^* and the $J/\psi + \pi \rightarrow D_1 \rightarrow D + D^*$ contribution. Of course, how big the cancellation actually is in this particular case has to be studied in detail and that will be done in the future work reported in [16].

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