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Hadronic form factors and the J/ψ secondary production cross section: An update

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Improving previous calculations, we compute the $D + \bar{D} \rightarrow J/\psi + \pi$ cross section using the most complete effective Lagrangians available. The new crucial ingredients are the form factors on the charm meson vertices, which are determined from QCD sum rules calculations. Some of them became available only very recently, and the last one, needed for our present purpose, is calculated in this work.

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I. INTRODUCTION

Before the BNL Relativistic Heavy Ion Collider (RHIC), relativistic heavy ion collision physics was relatively simple. We were basically searching for a quasi-ideal gas of deconfined quarks and gluons, called a quark-gluon plasma (QGP). One of the best signatures of this new state of matter was charmonium suppression [1]. During the last 4 years, intense work, both theoretical and experimental, has changed this naive picture drastically. On the theoretical side, careful numerical simulations [2–5] have shown that because of the importance of charm recombination in the deconfined phase and because of final state interactions, the number of J/ψ's may stay approximately the same. From the experimental side, especially from the analysis of elliptic flow, came the conclusion that the new state of matter is not a gas, but rather a strongly interacting fluid, now called sQGP [6]. Taking the existing calculations seriously, it is no longer clear that an overall suppression of the number of J/ψ's will be a signature of QGP. A more complex pattern can emerge, with suppression in some regions of the phase space and enhancement in others [7,8]. Whatever the new QGP signature (involving charm) turns out to be, it is necessary to understand better the mechanisms of J/ψ production and dissociation by collisions with comoving hadrons.

A great effort has been dedicated to understanding J/ψ dissociation in a hadronic environment. Since there is no direct experimental information on J/ψ absorption cross sections by hadrons, several theoretical approaches have been proposed to estimate their values. One approach was based on charm quark-antiquark dipoles interacting with the gluons of a larger (hadron target) dipole. This is the Bhanot-Peskin (BP) approach [9], which was rediscovered by Kharzeev and Satz [10] in the mid-1990s and updated [11,12] in the last few years. Finally, the recent next to leading order calculations presented in Ref. [13] have conclusively shown that, for charmonium, the formalism breaks down because this system is not heavy enough. Also considered was quark exchange between two (hadronic) bags [14,15]. The most explored approach has been the meson exchange mechanism [16–20]. In our opinion, the most reliable calculations of $\sigma_{J/\psi-\pi}$ were done with QCD sum rules [21]. However, due to a low momentum approximation, the validity of this calculation was restricted to low energy reactions, close to the dissociation threshold. This is probably not enough for the numerical simulations mentioned above. Therefore, to have cross sections valid at higher energies, the effective Lagrangrian approach still remains the most appropriate option.

After many works on the subject, some consensus has been achieved, at least regarding the determination of the order of magnitude, which, in the case of the J/ψ pion interaction, is determined to be $1 < \sigma_{J/\psi-\pi} < 10$ mb in the energy region close to the open charm production threshold.

Once the J/ψ dissociation cross section is known, using detailed balance one can attempt to estimate the charmonium formation cross section through the fusion of open charm, as, for example, $D\bar{D} \rightarrow J/\psi + \pi$. This is known as secondary charmonium production. As first pointed out in Ref. [22], in nucleus-nucleus collisions at sufficiently high energies, the number of produced D and D^* mesons increases, and the lifetime of the hadronic fireball increases. It becomes then possible that a significant number of J/ψ's be formed by open charm fusion. Later, an estimate made in Ref. [23] indicated that this mechanism would be relevant only for Large Hadron Collider (LHC) energies. The authors stressed, however, that their conclusion was very sensitive to the value of the J/ψ formation cross section, or equivalently to the absorption cross section, which in that case was the one computed with the Bhanot-Peskin approach. The subject was left aside for some time. Recently, after the revision of the J/ψ absorption cross section to larger values, secondary J/ψ production was incorporated into event generators in Refs. [24] and [5]. According to these simulations, the number of secondary J/ψ's is significant already at RHIC energies.

Given the renewed interest on the subject, we shall in this work further refine our estimate of the J/ψ interaction cross section, giving now emphasis to secondary charmonium production. We shall employ effective Lagrangians with form factors calculated with QCD sum rules (QCDSR). In particular, in the present calculation we shall make use of the D^*D^*J/ψ form factor, which was obtained only very recently [25], and we shall also calculate the $D^*D^*\pi$ form factor, which had not been calculated so far.

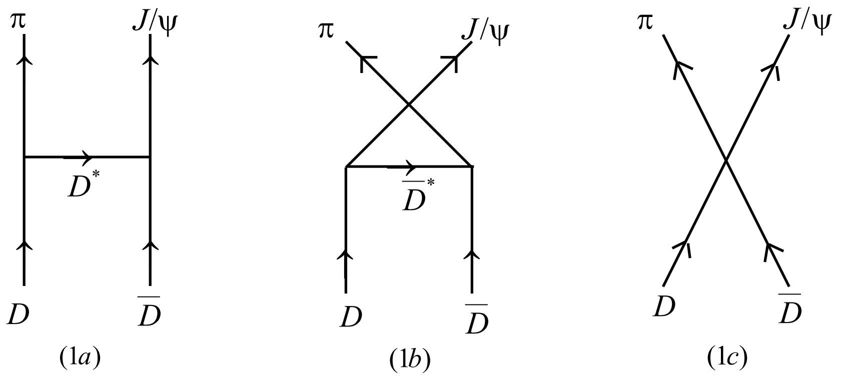

FIG. 1. Diagrams of contributions to the process $D\bar{D} \rightarrow J/\psi + \pi$.

II. THE EFFECTIVE LAGRANGIANS

Since the pioneering work of Matynian and Müller [16], there has been an intense discussion concerning the details and properties of the effective Lagrangians that describe the interactions among charm mesons. Here we do not add anything new to this discussion. We shall use what we believe is the state-of-the-art Lagrangian. For the sake of completeness and for future use, we present below the effective Lagrangians considered in this work:

$$\mathcal{L}_{D^*D\pi} = ig_{D^*D\pi}(D^*_\mu \partial^\mu \pi \bar{D} - D\partial^\mu \pi \bar{D}^*_\mu), \tag{1}$$

$$\mathcal{L}_{\psi D^* D} = g_{\psi D^* D}\epsilon^{\mu\nu\alpha\beta}\partial_\mu \psi_\nu (\partial_\alpha D^*_\beta \bar{D} + D\partial_\alpha \bar{D}^*_\beta), \tag{2}$$

$$\mathcal{L}_{\psi DD\pi} = ig_{\psi DD\pi}\epsilon^{\mu\nu\alpha\beta}\psi_\mu \partial_\nu D \partial_\alpha \pi \partial_\beta \bar{D}, \tag{3}$$

$$\mathcal{L}_{\psi DD} = -ig_{\psi DD}\psi^\mu(\partial_\mu D\bar{D} - D\partial_\mu \bar{D}), \tag{4}$$

$$\mathcal{L}_{D^*D^*\pi} = -g_{D^*D^*\pi}\epsilon^{\mu\nu\alpha\beta}\partial_\mu D^*_\nu \pi \partial_\alpha \bar{D}^*_\beta, \tag{5}$$

$$\begin{aligned}\mathcal{L}_{\psi D^* D^*} = {} & ig_{\psi D^* D^*}[\psi^\mu(\partial_\mu D^{*\nu}\bar{D}^*_\nu - D^{*\nu}\partial_\mu \bar{D}^*_\nu) \\ & + (\partial_\mu \psi_\nu D^{*\nu} - \psi_\nu \partial_\mu D^*_\nu)\bar{D}^{*\mu} + D^{*\mu} \\ & \times (\psi^\nu \partial_\mu \bar{D}^*_\nu - \partial_\mu \psi_\nu \bar{D}^{*\nu})],\end{aligned} \tag{6}$$

$$\mathcal{L}_{\psi D^* D\pi} = -g_{\psi D^* D\pi}\psi^\mu(D\pi \bar{D}^*_\mu + D^*_\mu \pi \bar{D}), \tag{7}$$

$$\begin{aligned}\mathcal{L}_{\psi D^* D^* \pi} = {} & ig_{\psi D^* D^* \pi}\epsilon^{\mu\nu\alpha\beta}\psi_\mu D^*{}_\nu \partial_\alpha \pi \bar{D}^*_\beta \\ & + ih_{\psi D^* D^* \pi}\epsilon^{\mu\nu\alpha\beta}\partial_\mu \psi_\nu D^*_\alpha \pi D^*{}_\beta.\end{aligned} \tag{8}$$

With these Lagrangians, we were are able to compute the process $D\bar{D} \rightarrow J/\psi + \pi$, which involves the diagrams in Fig. 1; the process $D^*\bar{D} \rightarrow J/\psi + \pi$, Fig. 2; and the process $D^*\bar{D}^* \rightarrow J/\psi + \pi$, Fig. 3.

Calling p_1 and p_2 the four-momenta of the incoming particles and p_3 and p_4 the four-momenta of the outgoing particles, we can derive the Feynman rules from the above Lagrangians and obtain the invariant amplitudes for each of the processes in Figs. 1–3. They are given by

$$\begin{aligned}\mathcal{M}^{1a}_\mu = {} & -g_{D^*D\pi}\, p_{3_\alpha} \frac{1}{t - m^2_{D^*}}\left(g_{\alpha\beta} - \frac{q_\alpha q_\beta}{m_{D^*}}\right) \\ & \times g_{\psi D^* D}\, \epsilon^{\rho\mu\theta\beta}\, q_\theta\, p_{4_\rho},\end{aligned} \tag{9}$$

$$\begin{aligned}\mathcal{M}^{1b}_\mu = {} & g_{D^*D\pi}\, p_{3_\alpha} \frac{1}{u - m^2_{D^*}} \\ & \times \left[g_{\alpha\beta} - \frac{(p_2 - p_3)_\alpha (p_2 - p_3)_\beta}{m_{D^*}}\right] \\ & \times g_{\psi D^* D}\, \epsilon^{\rho\mu\theta\beta}\, (p_2 - p_3)_\theta\, p_{4_\rho},\end{aligned} \tag{10}$$

$$\mathcal{M}^{1c}_\mu = g_{\psi D^* D\pi}\, \epsilon^{\mu\rho\theta\delta}\, p_{1_\rho}\, p_{2_\delta}\, p_{3_\theta}, \tag{11}$$

$$\mathcal{M}^{2a}_{\nu\sigma} = g_{D^*D\pi}\, p_{3_\nu} \frac{1}{t - m^2_D}\, g_{\psi DD}\, (p_2 - q)_\sigma, \tag{12}$$

$$\begin{aligned}\mathcal{M}^{2b}_{\nu\sigma} = {} & g_{\pi D^* D^*}\, \epsilon^{\gamma\nu\delta\alpha} p_{1_\gamma}\, q_\delta \frac{1}{t - m^2_{D^*}}\left(g_{\alpha\beta} - \frac{q_\alpha q_\beta}{m_{D^*}}\right) \\ & \times g_{\psi D^* D}\, \epsilon^{\rho\sigma\theta\beta} p_{4_\rho}\, q_\theta,\end{aligned} \tag{13}$$

$$\begin{aligned}\mathcal{M}^{2c}_{\nu\sigma} = {} & g_{D^*D\pi}\, p_{3_\alpha} \frac{1}{u - m^2_{D^*}} \\ & \times \left[g_{\alpha\beta} - \frac{(p_2 - p_3)_\alpha (p_2 - p_3)_\beta}{m_{D^*}}\right] \\ & \times g_{\psi D^* D^*}\, [(p_1 - p_2 + p_3)_\sigma\, g_{\nu\beta} - (p_1 + p_4)_\beta\, g_{\nu\sigma} \\ & + (p_2 - p_3 + p_4)_\nu\, g_{\sigma\beta}],\end{aligned} \tag{14}$$

$$\mathcal{M}^{2d}_{\nu\sigma} = -g_{\psi D^* D\pi}\, g_{\nu\sigma}, \tag{15}$$

$$\mathcal{M}^{3a}_{\mu\nu\sigma} = g_{D^*D\pi}\, p_{3_\mu} \frac{1}{t - m^2_D}\, g_{\psi D^* D}\, \epsilon^{\rho\sigma\theta\nu}\, p_{4_\rho}\, p_{2_\theta}, \tag{16}$$

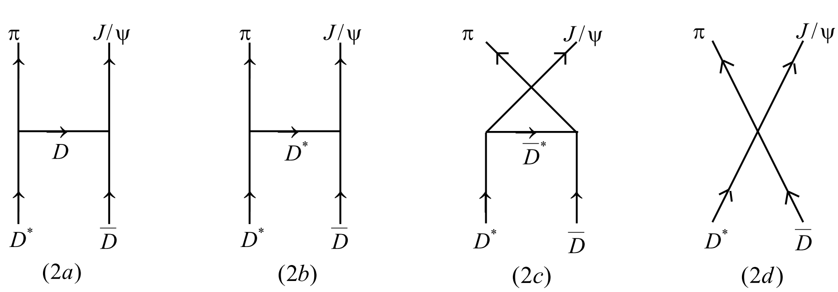

FIG. 2. Diagrams of contributions to the process $D^*\bar{D} \rightarrow J/\psi + \pi$.

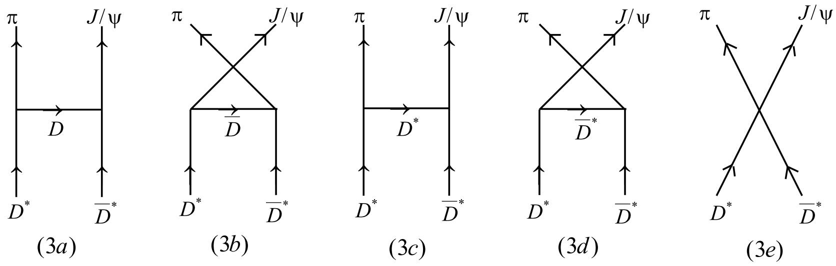

FIG. 3. Diagrams of contributions to the process $D^*\bar{D}^* \to J/\psi + \pi$.

$$\mathcal{M}^{3b}_{\mu\nu\sigma} = -g_{D^*D\pi}\, p_{3_\nu}\, \frac{1}{u - m_D^2}\, g_{\psi D^* D}\, \epsilon^{\rho\sigma\theta\mu}\, p_{4_\rho}\, p_{1_\theta}, \tag{17}$$

$$\begin{aligned}\mathcal{M}^{3c}_{\mu\nu\sigma} &= g_{\pi D^* D^*}\, \epsilon^{\delta\mu\theta\alpha}\, p_{1_\delta}\, q_\theta\, \frac{1}{t - m_{D^*}^2} \left(g_{\alpha\beta} - \frac{q_\alpha q_\beta}{m_{D^*}} \right) \\ &\quad \times g_{\psi D^* D^*}\, [(q - p_2)_\sigma\, g_{\nu\beta} + (p_2 + p_4)_\beta\, g_{\nu\sigma} \\ &\quad - (p_4 + q)_\nu\, g_{\sigma\beta}]\end{aligned} \tag{18}$$

$$\begin{aligned}\mathcal{M}^{3d}_{\mu\nu\sigma} &= -g_{\pi D^* D^*}\, \epsilon^{\delta\nu\theta\alpha}\, p_{2_\theta}\, (p_2 - p_3)_\delta\, \frac{1}{u - m_{D^*}^2} \\ &\quad \times \left[g_{\alpha\beta} - \frac{(p_2 - p_3)_\alpha (p_2 - p_3)_\beta}{m_{D^*}} \right] \\ &\quad \times g_{\psi D^* D^*}\, [(p_2 - p_3 - p_1)_\sigma\, g_{\mu\beta} + (p_1 + p_4)_\beta\, g_{\mu\sigma} \\ &\quad - (p_2 - p_3 + p_4)_\mu\, g_{\sigma\beta}],\end{aligned} \tag{19}$$

$$\mathcal{M}^{3e}_{\mu\nu\sigma} = -g_{\psi D^* D^* \pi}\, \epsilon^{\sigma\mu\theta\nu}\, p_{3_\theta} + h_{\psi D^* D^* \pi}\, \epsilon^{\sigma\mu\theta\nu}\, p_{4_\theta}. \tag{20}$$

Finally, the cross section for these processes is obtained with

$$\frac{d\sigma}{dt} = \frac{1}{N} \frac{1}{64\pi\, \mathbf{p}_i^2} \sum_{\text{spin}} |\mathcal{M}^2|, \tag{21}$$

where $\mathbf{p}_i^2$ a three-momentum squared in the center-of-mass system, and the factor $\frac{1}{N}$ comes from the average over the initial state polarizations.

As extensively discussed in previous works, although the above Lagrangians and amplitudes are quite satisfactory from the point of view of symmetry requirements, their straightforward application to the computation of cross sections leads to unacceptably large results. This comes from the fact that the exchanged particles may be far off-shell and therefore they enter (or leave) a vertex with a very different resolving power. In one extreme case considered in the recent past [26], a virtual J/ψ probing a D meson had the behavior of a parton (!). Of course, when this happens, the compact J/ψ almost misses the large D, and as a consequence the cross section of the whole process drops significantly. This physics of spatial extension and resolving power is contained in the form factors. Many authors have realized that calculations with and without form factors lead to results differing by up to two orders of magnitude! Therefore, we simply *cannot ignore the form factors*. We must include them in order to obtain reliable results!

Looking at the diagrams in Figs. 1–3, we notice that we need the following form factors (and the corresponding coupling constants):

$$\begin{matrix} g^{(D^*)}_{\pi D D^*}(t) & g^{(D)}_{\psi D D}(t) & g^{(D^*)}_{\psi D D^*}(t) \\ g^{(D)}_{\psi D D^*}(t) & g^{(D^*)}_{\psi D^* D^*}(t) & g^{(D^*)}_{\pi D^* D^*}(t), \end{matrix} \tag{22}$$

where t is the usual momentum transfer squared, and the superscript in parenthesis denotes the off-shell particle. This is an important distinction because the form factors in the same vertex are very different when different particles are off-shell. The most reliable way to compute these factors is with the use of the QCD sum rules techniques. The first one of the list was calculated in Ref. [27], the second in Ref. [26], and the third and fourth in Ref. [28], and the fifth in Ref. [25]; they read

$$g^{(D)}_{\pi D^* D}(t) = 17.9 \left[\frac{(3.5\,\text{GeV})^2 - m_D^2}{(3.5\,\text{GeV})^2 - t} \right] = h_4\big(t, m_D^2\big) \tag{23}$$

$$g^{(D)}_{\psi D D}(t) = 5.8\big[e^{-\left(\frac{20-t}{15.8}\right)^2}\big] = h_3(t) \tag{24}$$

$$g^{(D^*)}_{\psi D D^*}(t) = 20\big[e^{-\left(\frac{27-t}{18.6}\right)^2}\big] = h_1(t) \tag{25}$$

$$g^{(D)}_{\psi D D^*}(t) = 13\big[e^{-\left(\frac{26-t}{21.2}\right)^2}\big] = h_2(t) \tag{26}$$

$$g^{(\pi)}_{\pi D^* D^*}(t) = 4.8\big[e^{\left(\frac{t}{6.8}\right)}\big] = h_6(t, m_\pi^2) \tag{27}$$

$$g^{(D^*)}_{\psi D^* D^*}(t) = 6.2\big[e^{\left(\frac{t}{3.55}\right)}\big] = h_5(t) \tag{28}$$

The last form factor in Eq. (22) will be calculated below.

III. THE $\pi D^* D^*$ FORM FACTOR

In this section we shall, for the first time, compute the $\pi D^* D^*$ form factor using QCDSR [29,30]. In this approach, the short-range perturbative QCD is extended by an operator product expansion (OPE) expansion of the correlators, which results in a series in powers of the squared momentum with Wilson coefficients. The convergence at low momentum is improved by using a Borel transform. The expansion involves universal quark and gluon condensates. The quark-based calculation of a given correlator is equated to the same correlator calculated using hadronic degrees of freedom via a dispersion relation, thus providing sum rules from which a hadronic quantity can be estimated.

We shall use the three-point function to evaluate the $\pi D^* D^*$ form factor for an off-shell D^* meson, following the procedure suggested in Ref. [31] and further extended in Ref. [28]. This means that we shall calculate the correlators for a D^* off-shell and then for a π off-shell, requiring that the corresponding extrapolations to the respective poles lead to the same unique coupling constant.

The three-point function associated with a $\pi D^* D^*$ vertex with an off-shell D^* meson is given by

$$\begin{aligned}\Gamma^{(D^*)}_{\mu\nu}(p, p') &= \int d^4x\, d^4y\, \langle 0 | T\{ j_5(x)\, j_\nu(y)\, j^\dagger_\mu(0) \} | 0 \rangle \\ &\quad \times e^{ip'.x}\, e^{i(p-p').y},\end{aligned} \tag{29}$$

where $j_5 = i\bar{u}\gamma_5 d$, $j_\nu = \bar{c}\gamma_\nu u$, and $j_\mu = \bar{c}\gamma_\mu d$ are the interpolating fields for the π^-, D^{*0}, and D^{*-}, respectively, with u, d, and c being the up, down, and charm quark fields.

The phenomenological side of the vertex function, $\Gamma_{\mu\nu}(p, p')$, is obtained by considering π and D^* state contributions to the matrix element in Eq. (29):

$$\Gamma^{(\mathrm{phen})}_{\mu\nu}(p, p') = \frac{m_\pi^2 m_{D^*}^2}{m_u + m_d} \frac{F_\pi f_{D^*}^2 g^{(D^*)}_{D^* D^* \pi}(q^2)}{(q^2 - m_{D^*}^2)(p^2 - m_{D^*}^2)(p'^2 - m_\pi^2)} \times \varepsilon^{\mu\nu\lambda\delta} p_\lambda p'_\delta + \text{higher resonances}. \tag{30}$$

To derive Eq. (30), we made use of

$$\langle D^{*-}(p)|\pi^-(p')D^{*0}(q)\rangle = i g^{(D^*)}_{D^* D^* \pi}(q^2) \varepsilon^{\alpha\gamma\lambda\delta} p_\delta q_\lambda \epsilon_\alpha(p) \epsilon^*_\gamma(q), \tag{31}$$

where $q = p - p'$, and the decay constants F_π and f_{D^*} are defined by the matrix elements

$$\langle 0|j_5|\pi(p')\rangle = \frac{m_\pi^2 F_\pi}{m_u + m_d}, \tag{32}$$

and

$$\langle 0|j_\mu|D^*(p)\rangle = m_{D^*} f_{D^*} \epsilon_\mu(p), \tag{33}$$

where ϵ_μ is the polarization of the vector meson. The contribution of higher resonances and continuum in Eq. (30) will be taken into account as usual in the standard form of Ref. [32], through the continuun thresholds s_0 and u_0, for the D^* and π mesons, respectively.

The QCD or theoretical side of the vertex function is evaluated by performing Wilson's OPE of the operator in Eq. (29). Writing $\Gamma_{\mu\nu}$ in terms of the invariant amplitude,

$$\Gamma_{\mu\nu}(p, p') = \Lambda(p^2, p'^2, q^2) \varepsilon^{\mu\nu\lambda\delta} p_\lambda p'_\delta, \tag{34}$$

we can write a double dispersion relation for the invariant amplitude Λ over the virtualities p^2 and p'^2 holding $Q^2 = -q^2$ fixed

$$\Lambda^{(D^*)}(p^2, p'^2, Q^2) = -\frac{1}{4\pi^2} \int_{m_Q^2}^{s_0} ds \int_0^{u_0} du \frac{\rho(s, u, Q^2)}{(s - p^2)(u - p'^2)}, \tag{35}$$

where $\rho(s, u, Q^2)$ equals the double discontinuity of the amplitude $\Gamma(p^2, p'^2, Q^2)$ on the cuts $m_Q^2 \leqslant s \leqslant \infty$, $0 \leqslant u \leqslant \infty$, which can be evaluated using Cutkosky's rules. Finally, to suppress the condensates of higher dimension and at the same time reduce the influence of higher resonances, we perform a standard Borel transform [29]

$$\Pi(M^2) \equiv \lim_{n, Q^2 \to \infty} \frac{1}{n!} (Q^2)^{n+1} \left(-\frac{d}{dQ^2}\right)^n \Pi(Q^2), \tag{36}$$

(where $Q^2 = -q^2$ and the squared Borel mass scale $M^2 = Q^2/n$ is kept fixed) in both variables $P^2 = -p^2 \to M^2$ and $P'^2 = -p'^2 \to M'^2$ and equate the two representations described above. We get the following sum rule:

$$\frac{m_\pi^2 m_{D^*}^2}{m_u + m_d} F_\pi f_{D^*}^2 g^{(D^*)}_{\pi D^* D^*}(Q^2) e^{-m_\pi^2/M'^2} e^{-m_{D^*}^2/M^2} = (Q^2 + m_{D^*}^2) \bigg[\langle \bar{q}q \rangle \exp(-m_c^2/M^2) - \frac{1}{4\pi^2} \times \int_{m_c^2}^{s_0} ds \int_0^{u_{\max}} du \exp(-s/M^2) \exp(-u/M'^2) \times f(s, t, u) \theta(u_0 - u) \bigg], \tag{37}$$

where $t = -Q^2$,

$$f(s, t, u) = \frac{3 m_c u (2m_c^2 - s - t + u)}{[\lambda(s, u, t)]^{3/2}}, \tag{38}$$

$\lambda(s, u, t) = s^2 + u^2 + t^2 - 2su - 2st - 2tu$, and $u_{\max} = s + t - m_c^2 - st/m_c^2$.

We use the standard values for the numerical parameters: $m_{D^*} = 2.01$ GeV, $m_\pi = 140$ MeV, $F_\pi = \sqrt{2} \times 93$ MeV, $f_{D^*} = 240$ MeV, $m_u + m_d = 14$ MeV, $m_c = 1.3$ GeV, and $\langle \bar{q}q \rangle = -(0.23)^3$ GeV^3. For the continuum thresholds, we take $s_0 = (m_{D^*} + \Delta_s)^2$ with $\Delta_s = 0.5 \pm 0.1$ GeV and $u_0 = 1.4 \pm 0.2$ GeV^2.

In Fig. 4, we show the perturbative (dotted line) and the quark condensate (dashed line) contributions to the form factor $g^{(D^*)}_{\pi D^* D^*}(Q^2)$ at $Q^2 = 0.5$ GeV^2 as a function of the Borel mass M^2 at a fixed ratio $M'^2/M^2 = 0.64/(m_{D^*}^2 - m_c^2)$. We see that the quark condensate contribution is bigger than the perturbative contribution for values of the Borel mass smaller than $\sim$4.5 GeV^2. However, the sum of both contributions for the form factor is a very stable result as a function of the Borel mass. The quark condensate contribution grows with Q^2, while the perturbative contribution decreases. This imposes a limitation over the region of Q^2 that we can use to study the

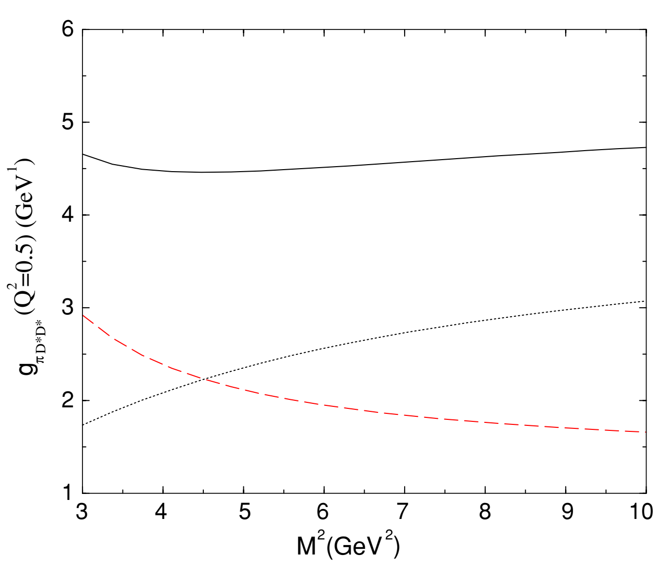

FIG. 4. (Color online) M^2 dependence of the perturbative contribution (dotted line) and the quark condensate contribution (dashed line) to the $g^{(D^*)}_{\pi D^* D^*}(Q^2)$ at $Q^2 = 0.5$ GeV^2. Solid line gives the final result for the form factor.

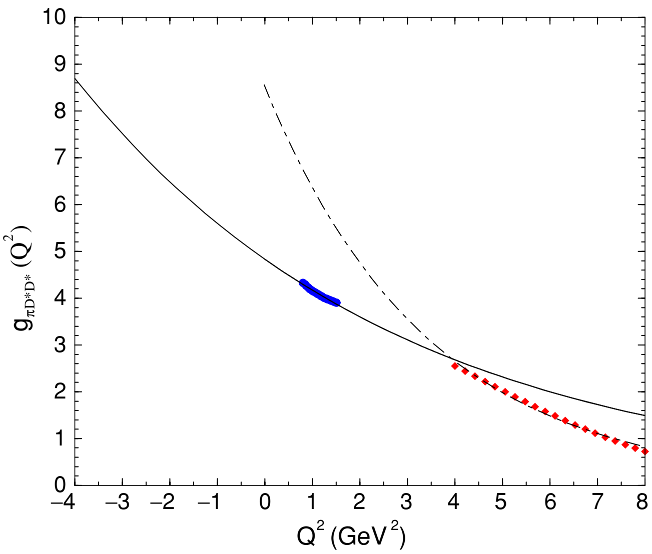

FIG. 5. (Color online) Momentum dependence of the $\pi D^* D^*$ form factors. Solid and dot-dashed lines give the parametrization of the QCDSR results for $g^{(D^*)}_{\pi D^* D^*}(Q^2)$ (circles) and $g^{(\pi)}_{\pi D^* D^*}(Q^2)$ (squares), respectively.

Q^2 dependence of the form factor. Fixing $M^2 = 10$ GeV^2, in Fig. 5 we show, through the filled circles, the momentum dependence of $g^{(D^*)}_{\pi D^* D^*}(Q^2)$.

Since the present approach cannot be used at $Q^2 < 0$, in order to extract the $g_{\pi D^* D^*}$ coupling from the form factor, we need to extrapolate the curve to the mass of the off-shell meson D^*. To do this, we fit the QCDSR results with an analytical expression. We have obtained a good fit using an exponential form

$$g^{(D^*)}_{\pi D^* D^*}(Q^2) = 4.8 e^{-Q^2/6.8}\,\text{GeV}^{-1}\,, \tag{39}$$

where 6.8 is in units of GeV^2. This fit is also shown in Fig. 5 through the solid line. From Eq. (39), we get $g_{\pi D^* D^*} = g^{(D^*)}_{\pi D^* D^*}(Q^2 = -m^2_{D^*}) = 8.7\,\text{GeV}^{-1}$. To check the consistency of the calculation, we also evaluate the form factor at the same vertex, but for an off-shell pion. In this case, we have to evaluate the three-point function

$$\Gamma^{(\pi)}_{\mu\nu}(p, p') = \int d^4x d^4y \langle 0|T\{j_\nu(x) j_5(y) j^\dagger_\mu(0)\}|0\rangle \times e^{ip'.x} e^{i(p-p')\cdot y}. \tag{40}$$

Proceeding in a similar way, we obtain the following sum rule:

$$\frac{m^2_\pi m^2_{D^*}}{m_u + m_d} F_\pi f^2_{D^*} g^{(\pi)}_{\pi D^* D^*}(Q^2) e^{-m^2_{D^*}/M'^2} e^{-m^2_{D^*}/M^2} = \left(Q^2 + m^2_\pi\right)\left[-\frac{1}{4\pi^2}\int_{s_{\min}}^{s_0} ds \int_{u_{\min}}^{u_0} du\, e^{-s/M^2} e^{-u/M'^2} \times \frac{3 m_c t\left(s + u - t - 2m^2_c\right)}{[\lambda(s, u, t)]^{3/2}}\right], \tag{41}$$

where $u_{\min} = m^2_c - \frac{m^2_c t}{s - m^2_c}$ and $s_{\min} = m^2_c(1 - \frac{t}{u_0 - m^2_c})$. Now we use $u_0 = (m_{D^*} + \Delta_u)^2$ with $\Delta_u = 0.5 \pm 0.1$ GeV, and $M^2 = M'^2$. The results are also rather stable as a function of the Borel mass. We also achieved a good fit of the QCDSR results for $g^{(\pi)}_{\pi D^* D^*}(Q^2)$ using an exponential form

$$g^{(\pi)}_{\pi D^* D^*}(Q^2) = 8.5\, e^{-Q^2/3.4}\ \text{GeV}^{-1}, \tag{42}$$

where 3.4 is in units of GeV^2. This fit is also shown in Fig. 5 through the dot-dashed line. From Eq. (42), we get $g_{\pi D^* D^*} = g^{(\pi)}_{\pi D^* D^*}(Q^2 = -m^2_\pi) = 8.5\ \text{GeV}^{-1}$.

Considering the uncertainties in the continuum thresholds and the difference in the values of the coupling extracted when the D^* or π mesons are off-shell, our result for the $\pi D^* D^*$ coupling constant is

$$g_{\pi D^* D^*} = 8.6 \pm 1.0\,\text{GeV}^{-1}. \tag{43}$$

The triple vertex couplings were calculated as explained above. The quartic vertex couplings could not be obtained with QCDSR, and we used the prescription given in Ref. [17],

$$g_{\psi DD\pi} = \left(\frac{\sqrt{3}}{6} - \frac{1}{4}\right)\frac{g_a N_c}{16\pi^2 F^3_\pi}, \tag{44}$$

$$g_{\psi D^* D\pi} = \frac{1}{2}\, g_{\psi DD}\, g_{D^* D\pi}, \tag{45}$$

$$g_{\psi D^* D^* \pi} = \frac{1}{2}\,\frac{g^3_a N_c}{32\pi^2 F_\pi}, \tag{46}$$

where g_a is obtained from

$$g_{\psi D^* D} = \frac{\sqrt{2}}{4\sqrt{3}}\,\frac{g^2_a N_c}{16\pi^2 F_\pi}. \tag{47}$$

In the above expressions, $N_c = 3$ and the triple vertex couplings are taken from our calculations. For completeness, we present in Table I all the couplings.

IV. THE CROSS SECTIONS

Having all the needed form factors, we now proceed to the evaluation of the cross sections. As in previous calculations, these cross sections for secondary J/ψ production will be related to the annihilation through the detailed balance

$$\sigma_{(3+4\to 1+2)} = \sigma_{(1+2\to 3+4)}\,\frac{(2S_1+1)(2S_2+1)}{(2S_3+1)(2S_4+1)}\,\frac{P^2_i}{P^2_f}. \tag{48}$$

TABLE I. Coupling constants used in the numerical calculations. First four were calculated with QCDSR; last three were obtained with the prescription of Ref. [17].

$g_{D^* D\pi}$	17.9
$g_{\psi D^* D}$	4.0 GeV^{-1}
$g_{\psi DD}$	5.8
$g_{\psi D^* D^*}$	6.2
$g_{D^* D^* \pi}$	8.6 GeV^{-1}
$g_{\psi DD\pi}$	10.0 GeV^{-3}
$g_{\psi D^* D\pi}$	51.9
$g_{\psi D^* D^* \pi}$	57.0 GeV^{-1}

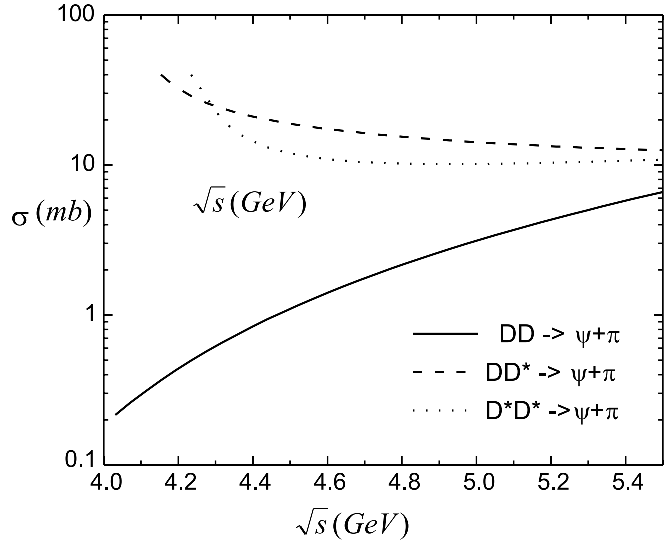

FIG. 6. J/ψ secondary production cross section without form factors.

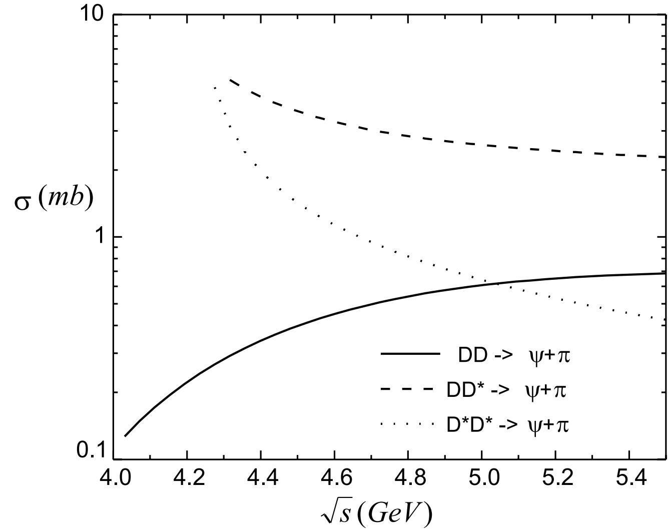

FIG. 8. J/ψ secondary production cross section with form factors.

In Fig. 6, we show the J/ψ secondary production cross section as a function of $\sqrt{s}$, without form factors. In all figures, the channels $D\bar{D} \rightarrow J/\psi + \pi$, $D\bar{D}^* \rightarrow J/\psi + \pi$, and $D^*\bar{D}^* \rightarrow J/\psi + \pi$ are represented by solid, dashed, and dotted lines, respectively. In Fig. 7, with the help of Eq. (48) we show the corresponding inverse reactions. As it can be seen, the cross sections have the same order of magnitude in both directions. Figures 8 and 9 are the analogs of 6 and 7 when we include the form factors in the calculations. Of course, as we stressed in the introduction, only these last two figures correspond to realistic numbers. The comparison of the two sets of figures is interesting only to estimate the effect of form factors. In previous studies doing the same kind of comparison, as for example in Ref. [17], the introduction of form factors reduced the cross sections by factors ranging between 20 and 50 depending on the channel. In that work, the form factor was the same for all vertices and the cutoff, not known, was estimated to be between 1 and 2 GeV. Our study is much more detailed and not only does each vertex has its own form factor, but also, depending on which particle is off-shell, the form factor is different. The final effect of all these peculiarities is the reduction of the cross sections by a factor around 7. Although significant, this reduction is smaller than previously expected.

Figure 8 contains our main results. The plotted cross sections can be compared with the results of Ref. [22] and, more directly, with Ref. [5]. In Fig. 2 of Ref. [22], although the variables in the plot are different, we can observe the same trend and relative importance of the three channels. In that work, the results were obtained with the quark model of Ref. [15]. Our curves share some features with the results of Ref. [5], such as, for example, the dominance of the DD^* channel and the falling trend of the DD^* and D^*D^* channels. The behavior of the DD channel is quite different. In the energy range of $\sqrt{s} > 4.5$ GeV, our cross sections are smaller by a factor of 2 (DD^*) or 5 (D^*D^* and DD).

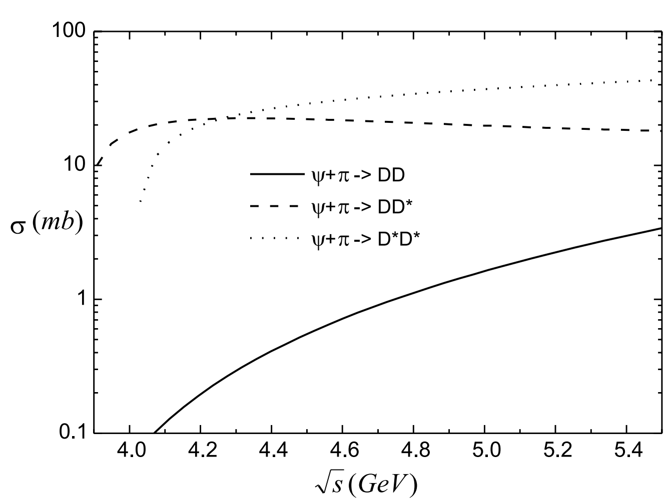

FIG. 7. J/ψ absorption cross section obtained through detailed balance without form factors.

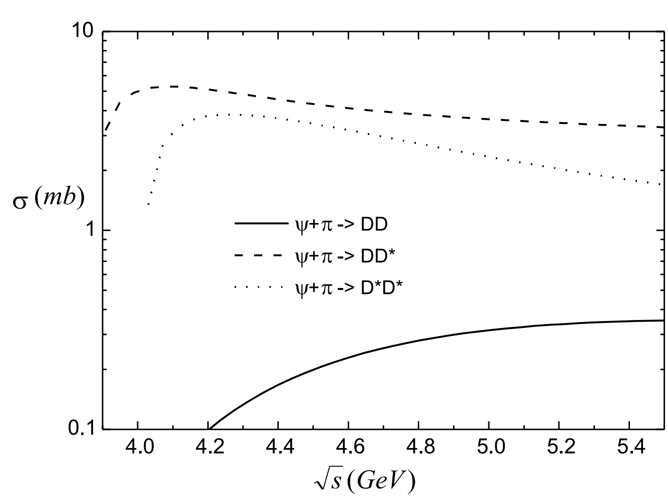

FIG. 9. J/ψ absorption cross section obtained through detailed balance with form factors.

These discrepancies are large but they are expected since in Ref. [5] all channels include the final state $J/\psi + \rho$. We could not include it consistently because the form factors of the $\rho D D^*$ and $\rho D^* D^*$ vertices have not yet been studied with our techniques and are thus not available. In the model used by the Giessen group, the cross sections for $D + \bar{D} \to J/\psi + \pi$ and $D + \bar{D} \to J/\psi + \rho$ are similar, and the same conclusion holds for the other inital state open charm mesons. If this would remain true in the effective Lagrangian approach, then our results including both final states would come closer to those of Ref. [5], thus giving more theoretical support to the model considered there.

V. SUMMARY AND CONCLUSIONS

We have updated the calculations of the cross sections for J/ψ dissociation and production in the effective Lagrangian approach. The novel feature introduced in this work is the use of the form factors (23)–(28) and especially (39), which was obtained here from QCDSR. We believe that our results are useful for numerical simulations of heavy ion collisions, such as those performed in Refs. [2–5] and [22–24]. The calculation of the cross sections of the processes considered here are complete. However, our program is not yet finished, and there are still form factors to be calculated, such as $\rho D D^*$ and $\rho D^* D^*$. These calculations are under way, and we will eventually have all hadronic form factors.

Although no strong statement can be made without knowing the $D + \bar{D} \to J/\psi + \rho$ cross section, our results give partial support to the conclusion advanced in Ref. [5], namely that the open charm fusion cross sections are large enough to produce a sizable number of “recreated” J/ψ’s already in heavy ion collisions at RHIC.

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