Eur. Phys. J. C 48, 545–551 (2006)
DOI 10.1140/epjc/s10052-006-0075-8

THE EUROPEAN PHYSICAL JOURNAL C

Regular Article – Theoretical Physics

J/ψ strong couplings to vector mesons

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Received: 10 May 2006 /
Published online: 24 October 2006 –

Abstract. We present a study of the cross sections $J/\psi X \to D^{(*)}\overline{D}^{(*)}$ ($X = \rho, \Phi$) based on the calculation of the effective tri- and four-linear couplings $J/\psi\, X\, D^{(*)}\, \overline{D}^{(*)}$ within a constituent quark model. In particular, the method for the calculation of the four-linear couplings is illustrated. The results obtained have been used in a recent analysis of J/ψ absorption by the hot hadron gas formed heavy-ion collisions at SPS energies.

PACS. 12.38.Mh; 12.39.-x; 25.75.-q

1 Introduction

The problem of computing J/ψ strong couplings to π, ρ and other pseudo-scalar and vector particles has its own interest because it opens the way to the calculation of cross sections of reactions like

$$J/\psi\, \{\pi, \rho, \dots\} \longrightarrow D^{(*)}\, \bar{D}^{(*)} . \tag{1}$$

Such cross sections are the basic ingredients to estimate the hadronic absorption background of J/ψ in heavy-ion collisions [1, 2]. The description of processes like those in (1) is a difficult task because no direct experimental information can be used and no first principle calculations are possible, so that one has to resort to building of models and making approximations to describe their dynamics.

The dissociation process of the J/ψ by hadrons has been considered in several studies [3–6], and the predicted cross sections show a different energy dependence and magnitude near threshold. Anyway, in all the models used in the literature, one consistently finds non-negligible cross section values (at least comparable with the nuclear one $\mathcal{N} J/\psi$, $\mathcal{N}$ = nucleon) especially for the reactions involving πs and ρs; for a review see for instance [7]. This is a clear indication that the picture of J/ψ absorption by nuclear matter, as an antagonist mechanism to plasma suppression, is incomplete as long as interactions with the hadron gas formed in nucleus–nucleus collisions (the so-called comoving particles) are not considered.

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The problem of calculating J/ψ dissociation by pseudoscalar and vector mesons has been addressed in [1, 2] within the *constituent-quark-meson* model (CQM), originally devised to compute exclusive heavy-light meson decays and tested on a quite large number of such processes [8–11]. The basic calculations refer to π and ρ contributions. The couplings to other mesons have been obtained under the hypothesis of flavour–octet symmetry.

The typical effective Feynman diagrams contributing to the J/ψ dissociation are depicted in Fig. 1. The trilinear couplings $\rho D^{(*)} \bar{D}^{(*)}$ have been calculated in [12] and the $J/\psi D^{(*)} \bar{D}^{(*)}$ couplings have been recently discussed in [13, 14] (we report these results at the end of Sect. 3), where also four-linear couplings involving pions have been derived. The aim of the present note is to explain the method used and the results obtained in evaluating the four-linear couplings of the kind $J/\psi \rho D^{(*)} \bar{D}^{(*)}$ (see Fig. 1, third diagram), not calculated elsewhere. The numerical values of the $J/\psi \Phi D_s^{(*)} \bar{D}_s^{(*)}$ couplings are also given. In the end, we present the cross section predictions, based on the complete set of contributing diagrams, for the processes $J/\psi\, \rho \to D^{(*)}\, \bar{D}^{(*)}$ and $J/\psi\, \Phi \to D_s^{(*)}\, \bar{D}_s^{(*)}$, together with an estimate of the associated theoretical uncertainties.

2 The model

In this section we summarize the basics of the CQM model (for a review see [8–11]) that will turn out to be useful to calculate the four-linear couplings for the reaction (see Fig. 2):

$$J/\psi\, \rho \longrightarrow D^{(*)}\, \bar{D}^{(*)} . \tag{2}$$

CQM is based on an effective Lagrangian which incorporates the heavy quark spin-flavor symmetries and the

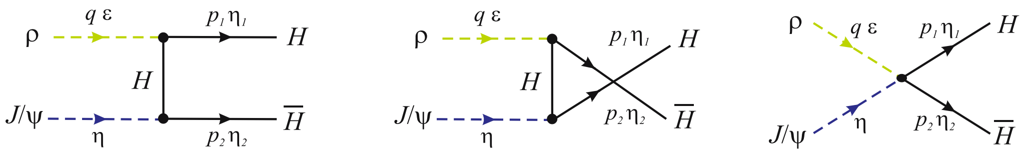

Fig. 1. Tree level effective Feynman diagrams for the $J/\psi\ \rho \to H\bar{H}$ reaction, H being $D^{(*)}$, with $D^{(*)} = D$ or D^*

chiral symmetry in the light sector. As in *heavy quark effective theory* (HQET) [15], we introduce the heavy superfield H describing the D and D^* charmed states, associated to the annihilation operators P_5, P^μ, respectively:

$$H(v) = \frac{1+\not{v}}{2}(\not{P} - P_5\gamma_5)\,, \tag{3}$$

where v is the four-velocity of the heavy meson; the limit of very large heavy quark mass is understood. The heavy quark propagator is

$$\frac{1+\not{v}}{2}\frac{\mathrm{i}}{v\cdot k}\,, \tag{4}$$

where k is the residual momentum defined by the equation $p_Q^\mu = m_Q v^\mu + k^\mu$ and related to the interaction of the light degrees of freedom with the heavy quark ($k \sim \mathcal{O}(\Lambda_{\rm QCD})$).

The CQM Lagrangian is obtained by applying bosonization techniques to Nambu–Jona–Lasinio (NJL) four fermion interactions of heavy and light quark fields [16, 17]; in particular, one can derive effective vertices between an heavy meson (H) and its constituent quarks (light quark (q) and heavy quark (Q_v)):

$$-\bar{q}\,\bar{H}\,Q_v + \mathrm{h.c.} \tag{5}$$

In Fig. 2 we show the typical diagrammatic equation to be solved in order to obtain g_4 (g_3), four- (tri-) linear couplings, in the various cases at hand (to obtain one of the relevant tri-linear couplings we could discard either the J/ψ or the ρ): the effective interaction at the meson level (l.h.s.) is modeled as an interaction at the quark–meson level (r.h.s.).

The vector mesons are introduced using a vector meson dominance (VMD) Ansatz: in the effective loop on the r.h.s. of Fig. 2 we have a vector current insertion on the heavy quark line c where the J/ψ is assumed to dominate the tower of 1^-, $c\bar{c}$ states mixing with the vector current [13, 14]). Similarly, ρ is coupled to the light quark current $q = (u, d)$.

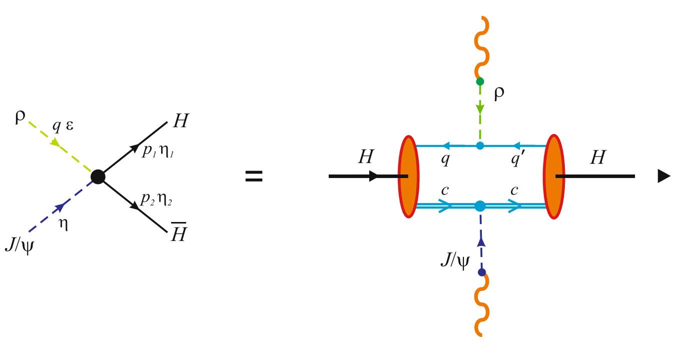

Fig. 2. Basic diagrammatic equation to compute the g_4 couplings. The l.h.s. is the effective vertex $J/\psi\rho H\bar{H}$ at meson level (ϵ and η are respectively the ρ and J/ψ polarizations), while the r.h.s. contains the 1-loop process to be calculated in the CQM model. On the light quark current and heavy quark current we have schematically represented the VMD mechanism

The ρ is described by the interpolating field ρ^μ and its kinetic term in the effective Lagrangian is built by the tensor field strength:

$$\mathcal{F}_{\mu\nu}(\rho) = \partial_\mu\rho_\nu - \partial_\nu\rho_\mu + [\rho_\mu, \rho_\nu]\,. \tag{6}$$

In the approach followed in [12], ρ^μ is defined by $\rho^\mu = im_\rho/f_\rho\hat{\rho}^\mu$ where $\hat{\rho}$ is a 3×3 hermitian traceless matrix analogous to the 3×3 π matrix of the pseudo-scalar SU_3 octet.

As in the Bando et al. approach [18] we introduce the ρ as a dynamical local gauge field of a hidden symmetry $SU(3)_V$ in the context of global $SU(3)_L \otimes SU(3)_R$ chiral symmetry. The ρ and J/ψ vertices with light and heavy quarks are given by

$$\bar{q}\left(\frac{m_\rho^2}{f_\rho}\,\rho^\mu\gamma_\mu\right)q \tag{7}$$

and

$$\bar{Q}_v\left(\frac{m_J^2}{f_J\,m_J}\,J^\mu\gamma_\mu\right)Q_v\,, \tag{8}$$

respectively. The decay constants f_ρ and f_J are defined by

$$\langle 0|V_\mu|\rho\rangle = \mathrm{i}f_\rho\,\epsilon_\mu\,,\quad \langle 0|V_\mu|J/\psi\rangle = \mathrm{i}m_J\,f_J\,\eta_\mu\,, \tag{9}$$

where $f_\rho = 0.152\,\mathrm{GeV}^2$ and $f_J = 0.405\,\mathrm{GeV}$.

Once having established the form of the effective vertices occurring in the loop diagram in the r.h.s. of Fig. 2, one has to compute it using some regularization; we will adopt the Schwinger proper time.

CQM does not include any confining potential and an infrared cutoff μ is needed to prevent low integration momenta to access the energy region where confinement should be at work. The kinematic condition for the free dissociation of H in heavy (m_Q) quark and light (m) quark is given by

$$m_H > m_Q + m\,, \tag{10}$$

with $p_H = m_H\,v \approx m_Q\,v + k$. It follows that

$$k\cdot v > m\,. \tag{11}$$

In the hadron rest frame we have $k^0 > m$ and we can therefore require

$$\mu \approx m\,. \tag{12}$$

The value of the constituent light quark mass in the model at hand is given by a gap equation [8–11]:

$$m - m_0 - 8\, G\, I_1(m^2) = 0\,, \tag{13}$$

where $G = 5.25\,\text{GeV}^{-2}$, m_0 is the current mass, and the I_1 integral is defined in the appendix. As a consistency check, putting a zero current mass for the u, d species we get a constituent mass of 300 MeV, while for a strange current mass of $m_0 = 131$ MeV we obtain a strange constituent mass of 500 MeV using $\mu = 300, 500$ MeV respectively in the calculation of I_1.

The residual momentum has an upper limit given by the chiral symmetry braking scale $\Lambda_\chi \simeq 4\pi f_\pi$ which we adopt as a UV cutoff [8–11].

The momenta running in the loop are therefore limited by two cutoffs: μ and Λ. These two cutoffs are implemented by the Schwinger regularization on the light propagator as follows:

$$\int \mathrm{d}^4 l \frac{1}{(l^2 - m^2)} \longrightarrow \int \mathrm{d}^4 l \int_{1/\Lambda^2}^{1/\mu^2} \mathrm{d}s\, \mathrm{e}^{-s\left(l^2 + m^2\right)}\,. \tag{14}$$

The diagrammatic equation in Fig. 2 states that the effective vertex $J/\psi\, \rho\, H\, \bar{H}$ is given by

$$\begin{aligned} \langle H|\, \mathrm{i}\mathcal{L}_{\rho\, J\, H\, H}\, |H\, \rho\, J/\psi\rangle \;&=\; (-1)\sqrt{Z_H m_H Z_{H'} m_{H'}} \\ \times N_c \int \frac{\mathrm{d}^4 l}{(2\pi)^4} \mathrm{Tr}\Bigg[&\left(-\mathrm{i}\bar{H}'(v')\right) \frac{\mathrm{i}}{v'\cdot l + \Delta}\, \frac{m_J^2}{f_J\, m_J} \not\eta\, \frac{\mathrm{i}}{v\cdot l + \Delta} \\ \times \left(-\mathrm{i}H(v)\right)\, &\frac{\mathrm{i}}{\not l - m}\, \mathrm{i}\frac{m_\rho^2}{f_\rho} \not\epsilon\, \frac{\mathrm{i}}{\not l + \not q - m} \Bigg]\,. \end{aligned} \tag{15}$$

H and $\bar{H}'$ represent the heavy-light external meson fields, given in (3), labeled by their four-velocities v, v', while the $\sqrt{Z_H m_H Z_{H'} m_{H'}}$ coupling factor of heavy mesons to quarks is calculated in [8–11]. m is the constituent mass of a light quark $q = u, d$. $\mathcal{L}_{\rho\, J\, H\, H}$ is the effective Lagrangian that contains the effective couplings g_4. The parameter Δ appearing in the heavy propagator is defined by

$$\Delta = M_H - m_Q\,, \tag{16}$$

i.e., the mass of the heavy-light meson minus the mass of the heavy quark contained in it. Δ is the main free parameter of the model. It varies in the range $\Delta = 0.3, 0.4, 0.5$ GeV for u, d light quarks and $0.5, 0.6, 0.7$ GeV for strange quarks [20]. Varying Δ gives an estimate of the theoretical error.

3 The calculation

The ρ is coupled to the light quarks by VMD, ϵ being its polarization and q its 4-momentum. The J/ψ, having polarization η, is also coupled via VMD, but to the heavy quarks (η appears in the trace between the two heavy quark propagators, while ϵ appears between the two light quark propagators). In front of this expression we have the fermion loop factor. Using the Feynman trick the fermion loop on the r.h.s. of (15) becomes

$$\begin{aligned} &\frac{m_J\, m_\rho^2}{f_J\, f_\rho} \sqrt{Z_H m_H Z_{H'} m_{H'}} \int_0^1 \mathrm{d}x \frac{\partial}{\partial \tilde{m}^2}\, iN_c \\ &\int \frac{\mathrm{d}^4 l}{(2\pi)^4} \frac{\mathrm{Tr}\left[\bar{H}'\, \not\eta\, H\, (\not l - \not q x + m)\, \not\epsilon\, (\not l - \not q x + \not q + m)\right]}{(l^2 - \tilde{m}^2)\, (v\cdot l + \delta)\, (v'\cdot l + \delta')}\,, \end{aligned} \tag{17}$$

in which we have defined

$$\tilde{m}^2 = m^2 + x\, m_\rho^2\, (x - 1)\,, \tag{18}$$

with $m = 0.3$ GeV the constituent mass for a light quark, u, d, and

$$\delta = \Delta - x\, q\cdot v \approx \Delta - x\, E_\rho\,, \tag{19}$$

$$\delta' = \Delta - x\, q\cdot v' \approx \Delta - x\, \omega\, E_\rho\,, \tag{20}$$

where E_ρ is the energy of the incident ρ and $\omega = vv'$ ($v' = \omega v$). The cross section computation is performed in the frame where J/ψ is at rest. ω, in this frame, is related to the meson masses by

$$\omega = \frac{m_{J/\psi}^2 + m_\rho^2 - m_H^2 - m_{H'}^2 + 2E_\rho m_{J/\psi}}{2m_H m_{H'}}\,. \tag{21}$$

By kinematic considerations the energy threshold of the reactions (2) for the $D\bar{D}$ and $D^*\bar{D}$ channels is $E_\rho \simeq 0.77$ GeV, whereas for the $D^*\bar{D}^*$ channel it is $E_\rho \simeq 0.96$ GeV, with $\omega \approx 1$. We consider ρ particles with energies in the range between 0.77 and 1 GeV where the two final state mesons are almost at rest.

The trace computation in (17) will introduce a number of scalar combinations of the momenta and polarizations of the external particles. Each of these combinations will be weighted by a scalar integral which amounts to a numerical factor: the couplings. All the couplings that we can extract by direct computation can be written in terms of five basic expressions: L_5, A, B, C, D. The latter are linear combinations of the I_i, L_i integrals listed in the appendix and are defined by

$$\frac{\partial}{\partial \tilde{m}^2} \mathrm{i}N_c \int \frac{\mathrm{d}^4 l}{(2\pi)^4} \frac{1}{(l^2 - \tilde{m}^2)\, (v\cdot l + \delta)\, (v'\cdot l + \delta')} = L_5\,, \tag{22}$$

$$\begin{aligned} &\frac{\partial}{\partial \tilde{m}^2} \mathrm{i}N_c \int \frac{\mathrm{d}^4 l}{(2\pi)^4} \frac{l_\mu}{(l^2 - \tilde{m}^2)\, (v\cdot l + \delta)\, (v'\cdot l + \delta')} \\ &= A\left(v_\mu + v'_\mu\right)\,, \end{aligned} \tag{23}$$

$$\begin{aligned} &\frac{\partial}{\partial \tilde{m}^2} \mathrm{i}N_c \int \frac{\mathrm{d}^4 l}{(2\pi)^4} \frac{l_\mu l_\nu}{(l^2 - \tilde{m}^2)\, (v\cdot l + \delta)\, (v'\cdot l + \delta')} \\ &= B g_{\mu\nu} + C\left(v_\mu v_\nu + v'_\mu v'_\nu\right) + D\left(v_\mu v'_\nu + v'_\mu v_\nu\right)\,. \end{aligned} \tag{24}$$

After performing the loop integrals we write the final form for the effective Lagrangian $\mathcal{L}_{\rho\, J\, H\, H}$ in (15) for the four-linear coupling:

$$\begin{aligned}\mathcal{L}_{\rho\, J\, D\, D} = J_\alpha\, \partial_\mu \rho_\nu \big[& g_1(\partial^\mu D\, \partial^\nu \partial^\alpha \bar{D} - \partial^\nu D\, \partial^\mu \partial^\alpha \bar{D}) \\ & + g_2(\partial^\mu \bar{D}\, \partial^\nu \partial^\alpha D - \partial^\nu \bar{D}\, \partial^\mu \partial^\alpha D)\big] \\ & + \partial_\mu \rho_\nu \big[g_3(J^\mu\, \partial^\nu \bar{D} - J^\nu\, \partial^\mu \bar{D}) D \\ & + g_4(J^\mu\, \partial^\nu D - J^\nu\, \partial^\mu D)\bar{D}\big] \\ & + g_5 J_\mu \rho_\nu [\partial^\mu \partial^\nu \bar{D}\, D + \partial^\mu \partial^\nu D\, \bar{D}] \\ & + g_6 \big[J \cdot \partial D\, \rho \cdot \partial \bar{D} + J \cdot \partial \bar{D}\, \rho \cdot \partial D \big] \\ & + g_7 J \cdot \rho\, D\, \bar{D}\,, \end{aligned} \tag{25}$$

where the velocities of the heavy mesons are converted to derivatives of the corresponding fields.

Similarly

$$\begin{aligned}\mathcal{L}_{\rho\, J\, D^*\, D} = \varepsilon_{\alpha\beta\gamma\delta} \Bigg\{ & - h_1\, \partial^\alpha \rho^\beta J^\gamma \bar{D}^{*\delta} D \\ & - h_7 \frac{\partial^\alpha \rho^\mu J^\beta \partial^\gamma \bar{D}^*{}_\mu \partial^\delta D}{m_{D^*} m_D} \\ & - h_4 \frac{\partial^\alpha \rho^\beta \partial^\mu \bar{D}^{*\gamma} \partial^\delta D J_\mu}{m_{D^*} m_D} \\ & + h_2 \Bigg[\left(\frac{\partial^\alpha \rho^\beta J^\mu}{m_D^2} + \frac{\partial^\alpha \rho^\mu J^\beta}{m_D^2} - \frac{\partial^\mu \rho^\alpha J^\beta}{m_D^2} \right) \\ & \times \bar{D}^{*\gamma}\, \partial^\delta \partial_\mu D - \frac{\partial^\alpha \rho^\beta J^\gamma}{m_D^2}\, \bar{D}^{*\mu}\, \partial^\delta \partial_\mu D \\ & - \partial^\mu \rho^\alpha \left(J^\beta \frac{\partial_\mu \bar{D}^{*\gamma}}{m_{D^*} m_D} + \frac{\partial^\gamma \bar{D}^{*\beta} J_\mu}{m_{D^*} m_D} \right) \partial^\delta D \\ & + \left(\frac{\partial^\alpha \rho^\mu J^\beta \partial^\delta \bar{D}^{*\gamma}}{m_{D^*} m_D} - \frac{\partial^\alpha \rho^\beta J^\gamma \partial^\delta \bar{D}^{*\mu}}{m_{D^*} m_D} \right) \partial_\mu D \Bigg] \\ & + h_5 \Bigg[J^\beta \left(\frac{\partial^\alpha \rho^\mu - \partial^\mu \rho^\alpha}{m_{D^*} m_D} \right) \partial^\delta \bar{D}^{*\gamma}\, \partial_\mu D \\ & - \frac{\partial^\alpha \rho^\beta\, \partial^\gamma \bar{D}^{*\mu} \partial^\delta D\, J_\mu}{m_{D^*} m_D} \\ & + \left(\frac{\partial^\alpha \rho^\mu J^\beta}{m_{D^*}^2} - \frac{\partial^\alpha \rho^\beta J_\mu}{m_{D^*}^2} - \frac{\partial^\mu \rho^\alpha J^\beta}{m_{D^*}^2} \right) \\ & \times \partial^\delta \partial_\mu \bar{D}^{*\gamma} D \Bigg] \\ & + h_6 \Bigg[J^\beta \\ & \times \left(\frac{\partial^\mu \rho^\alpha \partial^\delta \bar{D}^{*\gamma} \partial_\mu D}{m_{D^*} m_D} - \frac{\partial^\alpha \rho^\mu \partial_\mu \bar{D}^{*\gamma} \partial^\delta D}{m_{D^*} m_D} \right) \\ & + \partial^\mu \rho^\alpha \left(\frac{J^\beta\, \partial^\gamma \bar{D}^*{}_\mu}{m_{D^*} m_D} - \frac{\partial^\gamma \bar{D}^{*\beta}\, J_\mu}{m_{D^*} m_D} \right) \partial^\delta D \\ & + \left(\frac{\partial^\mu \rho^\alpha J^\beta}{m_{D^*}^2} - \frac{\partial^\alpha \rho^\beta J^\mu}{m_{D^*}^2} - \frac{\partial^\alpha \rho^\mu J^\beta}{m_{D^*}^2} \right) \\ & \times \partial^\delta \partial_\mu \bar{D}^{*\gamma}\, D \Bigg] + h_3 \\ & \times \left(\frac{\partial^\alpha \rho^\beta \partial^\delta \bar{D}^{*\gamma} J_\mu\, \partial^\mu D - \partial^\mu \rho^\alpha J^\beta \partial^\gamma \bar{D}^*{}_\mu \partial^\delta D}{m_{D^*} m_D} \right) \\ & + h_8\, \rho^\alpha J^\beta \left(\frac{\bar{D}^{*\gamma} \partial^\delta D}{m_D} - \frac{\partial^\delta \bar{D}^{*\gamma} D}{m_{D^*}} \right) \\ & - 2h_9 \frac{J^\alpha \partial^\gamma \partial^\mu \bar{D}^{*\beta} \partial^\delta D \rho_\mu}{m_{D^*}^2 m_D} \\ & - 2h_{10} \frac{J^\alpha \partial^\gamma \bar{D}^{*\beta} \partial^\delta \partial_\mu D \rho^\mu}{m_{D^*} m_D^2} \Bigg\} + \text{h.c.} \end{aligned} \tag{26}$$

We omit here the explicit expression for the lagrangian $\mathcal{L}_{\rho\, J\, D^*\, D^*}$. In Table 1 we list the expressions of the couplings and their numerical values with theoretical errors; we call them $g_4 = \{g, h, f\}$, related to the four-linear couplings $J/\psi\, \rho\, D\, \bar{D}$, $J/\psi\, \rho\, D^*\, \bar{D}$, $J/\psi\, \rho\, D^*\, \bar{D}^*$ respectively. Our central values for the couplings are obtained for $\Delta = 0.4\,\text{GeV}$ (and $Z_H = 2.36\,\text{GeV}^{-1}$) [8–11], while for m_H we use the experimental value for the $D^{(*)}$ mass.

Couplings are typically complicated functions of E_ρ and it is not as easy as in [13, 14] to determine an explicit polar form factor dependence common to all of them. On the other hand we have in mind to adopt these couplings to compute cross sections of the kind $\sigma_{J/\psi\rho \to \text{open charm}}$ and use this information to compute thermal averages of the form

$$\left\langle \rho \cdot \sigma_{J/\psi\rho \to \text{open charm}} \right\rangle_T = \frac{N_c}{2\pi^2} \int_{E_0}^{\infty} \mathrm{d}E_\rho\, \frac{p_\rho\, E_\rho\, \sigma(E_\rho)}{\mathrm{e}^{E_\rho/kT} - 1}\,, \tag{27}$$

where E_0 is the energy threshold needed to open the reaction channel and the Bose distribution is used to describe an ideal gas of mesons. N_c, the number of colors, is set equal to 3 in all our calculations. ρ in the l.h.s. of (27) is the number density of particles in the gas. The Boltzmann factor $\exp(-E_\rho/T)$ will be at work as an exponential form factor cutting high energy tails faster than any polar one. We could therefore avoid any arbitrary Ansatz on form factors at the interaction vertices. We limit ourself to the study of the dependence of our couplings on E_ρ in the range of energy where we reasonably think to have ρ mesons in the hadron gas excited by a peripheral heavy-ion collision. Estimating the J/ψ absorption background to the suppression signal in heavy-ion collisions amounts to compute the attenuation lengths (inverse of the thermal averages in (27)) in a hot gas, $T \approx 170\,\text{MeV}$, populated by $\pi, K, \eta, \rho, \omega, \ldots$ mesons. We could therefore expect $E_\rho \approx 770$–$1000\,\text{MeV}$.

We have considered also the interaction of J/ψ with Φ. The reaction in this case is

$$J/\psi\, \Phi \longrightarrow D_s^{(*)}\, \bar{D}_s^{(*)}\,. \tag{28}$$

The loop that must be calculated in the case in which we substitute a Φ particle to the ρ is the same as in Fig. 2 but with $q, q' = s$ ($m_s = 500\,\text{MeV}$, $\mu \sim m_s$), and with superfields H_s in place of H.

The structure of the coupling of Φ to the light quark current is identical to the ρ one, and all the above equations are valid also in this case with the substitutions $m_\rho \to m_\Phi$, $f_\rho \to f_\Phi$, $H \to H_s$. Numerically we have used $f_\Phi = 0.249\,\text{GeV}^2$, while for m_{H_s} we have used the experimental value for $D_s^{(*)}$. The numerical values of the four-linear coupling $J/\psi\, \Phi\, H\, \bar{H}$ are reported in Table 1.

Table 1. The four-linear couplings $g_4 = \{g_i, h_i, f_i\}$ expressed as linear combinations of the basic scalar integrals listed in the appendix. All expression must be understood as $\frac{m_J\, m_\rho^2}{f_J\, f_\rho}\sqrt{Z_H m_H Z_{H'} m_{H'}}\int_0^1 \mathrm{d}x\, g_4$: the mean values are estimated by setting $\Delta = 0.4$ GeV ($\Delta = 0.6$ GeV) and varying the energy of ρ (Φ) in the range $E_\rho = 0.770$–1 GeV ($E_\Phi = 1.02$–1.2 GeV), while the error is calculated by combining the excursion in the selected energy range at $\Delta = 0.4$ GeV with the uncertainty over Δ (in the range $\Delta = 0.5$–0.7 GeV). Some of the couplings are unavoidably affected by large uncertainties. This affects the determination of the attenuation lengths discussed in the text to the extent pointed out in [1, 2]

	$J/\psi X D\bar{D}$	$X=\rho$	$X=\Phi$	
g_1	$\frac{2}{m_D^3} Ax$	2 ± 1	0.7 ± 0.2	[GeV^{-4}]
g_2	$\frac{2}{m_D^3} A(x-1)$	-3.7 ± 0.7	-1.6 ± 0.2	[GeV^{-4}]
g_3	$\frac{1}{m_D}\left(2A(1-x+\omega x) - mL_5\right)$	25 ± 4	12 ± 1	[GeV^{-2}]
g_4	$\frac{1}{m_D}\left(2A(\omega+x+\omega x) - mL_5\right)$	43 ± 10	20 ± 3	[GeV^{-2}]
g_5	$\frac{2}{m_D^2}(C+D-mA)$	-8 ± 3	-6 ± 1	[GeV^{-2}]
g_6	$-\frac{1}{m_D^2}\left(2(mA+B+D(\omega-1)) + \left(m_\rho^2 x(x-1) - m^2\right) L_5\right)$	-15 ± 6	-12 ± 2	[GeV^{-2}]
g_7	$\left(2B+2C+2\omega D + \left(m_\rho^2 x(x-1) - m^2\right) L_5\right)(\omega-1)$	-0.4 ± 0.4	-0.4 ± 0.2	
	$J/\psi X D^*\bar{D}$	$X=\rho$	$X=\Phi$	
h_1	$(mL_5 + (A-B)x)(\omega-1)$	1 ± 2	0.1 ± 0.6	[GeV^{-1}]
h_2	$B(x-1)$	-9 ± 4	-5 ± 1	[GeV^{-1}]
h_3	$mL_5 - Bx$	-6 ± 12	-6 ± 3	[GeV^{-1}]
h_4	$mL_5 + A(x-1) - B$	-35 ± 15	-20 ± 4	[GeV^{-1}]
h_5	A	35 ± 11	15.7 ± 3	[GeV^{-1}]
h_6	Ax	15 ± 8	6 ± 2	[GeV^{-1}]
h_7	$B - mL_5$	16 ± 16	11 ± 4	[GeV^{-1}]
h_8	$\left(m^2 + m_\rho^2 x(1-x)\right) L_5 - 2C - D - E - 2F\omega$	1.3 ± 2	1.1 ± 0.8	
h_9	$D + F - mA$	-19 ± 7	-17 ± 3	
h_{10}	$E + F - mB$	-15 ± 6	-13 ± 2	
	$J/\psi X D^*\bar{D}^*$	$X=\rho$	$X=\Phi$	
f_1	$\frac{1}{m_{D^*}}(A+B-mL_5)$	25.5 ± 0.6	13.6 ± 0.1	[GeV^{-2}]
f_2	$\frac{1}{m_{D^*}}(B - mL_5 + A(2\omega x - 2x + 1))$	26 ± 1	13.86 ± 0.08	[GeV^{-2}]
f_3	$\frac{2}{m_{D^*}^3} Ax$	4 ± 2	1.5 ± 0.5	[GeV^{-4}]
f_4	$\frac{1}{m_{D^*}}(A - mL_5 + B(-2\omega x + 2x + 2\omega - 1))$	26.0 ± 0.5	13.8 ± 0.1	[GeV^{-2}]
f_5	$\frac{2}{m_{D^*}^3} B(x-1)$	-2.2 ± 0.7	-1.2 ± 0.1	[GeV^{-4}]
f_6	$\left(-m^2 L_5 + 2C + D + E + 2F\omega + m_\rho^2(x^2-x)L_5\right)(1-\omega)$	0.03 ± 0.1	0.03 ± 0.03	
f_7	$\frac{1}{m_{D^*}^2}\left(-m^2 L_5 + 2C + D + E + 2F\omega + m_\rho^2(x^2-x)L_5\right)$	-0.1 ± 2	-0.2 ± 0.2	[GeV^{-2}]
f_8	$\frac{2}{m_{D^*}^2}(D + F - mA)$	-8 ± 2	-8.2 ± 0.3	[GeV^{-2}]
f_9	$\frac{1}{m_{D^*}^2}\left(-m^2 L_5 + 2mA + 2C - D + E + 2F(\omega-1) + m_\rho^2(x^2-x)L_5\right)$	8.3 ± 2	8.0 ± 0.1	[GeV^{-2}]
f_{10}	$\frac{1}{m_{D^*}^2}\left(-m^2 L_5 + 2mB + 2C + D - E + 2F(\omega-1) + m_\rho^2(x^2-x)L_5\right)$	7.3 ± 0.8	6.3 ± 0.2	[GeV^{-2}]
f_{11}	$\frac{2}{m_{D^*}^2}(E + F - mB)$	-7 ± 1	-6.5 ± 0.5	

For completeness we conclude this section including the tri-linear couplings ρHH and $J/\psi HH$; ρHH are computable within the same framework by simply not including the J/ψ in the diagrammatic equation of Fig. 2. The vertex $\rho D^{(*)}\bar{D}^{(*)}$ is described by the two constants $g_3 = \beta, \lambda$ and the effective Lagrangian describing this interaction can be written as [8–12]

$$\mathcal{L}_{HH\rho} = -\mathrm{i}\beta\, \mathrm{Tr}[\bar{H}H]\, v\cdot\rho + \mathrm{i}\lambda\, \mathrm{Tr}\left[\bar{H}\sigma_{\mu\nu}H\right]\mathcal{F}^{\mu\nu}\,, \quad (29)$$

where the field ρ^μ and the tensor $\mathcal{F}^{\mu\nu}$ have been defined above. The numerical values are

$$\beta = -0.98\,, \quad \lambda = +0.42\ \mathrm{GeV}^{-1}\,; \quad (30)$$

analogously β_Φ and λ_Φ for the $\mathcal{L}_{HH\Phi}$ are

$$\beta_\Phi = -0.48\,, \quad \lambda_\Phi = +0.14\ \mathrm{GeV}^{-1}\,. \quad (31)$$

As for the couplings $J/\psi D^{(*)}\bar{D}^{(*)}$, they have been extensively discussed in [13, 14]. Here we just report the main results. Observe that the effective Lagrangian in this case is

$$\mathcal{L}_{JHH} = \mathrm{i}g_{JHH}\,\mathrm{Tr}[\bar{H}\gamma_\mu H]J^\mu\,, \tag{32}$$

where H can be any of the pairs $D\,D^*$ or $D_s\,D_s^*$ (neglecting $SU(3)$ breaking effects). As a consequence of the spin symmetry of the HQET we find

$$g_{JD^*D^*} = g_{JDD}\,, \quad g_{JDD^*} = \frac{g_{JDD}}{m_D}\,. \tag{33}$$

The numerical values are given by

$$\begin{aligned} g_{JDD} &= 8.0 \pm 0.5\,,\\ g_{JDD^*} &= 4.05 \pm 0.25\,\mathrm{GeV}^{-1}\,,\\ g_{JD^*D^*} &= 8.0 \pm 0.5\,. \end{aligned} \tag{34}$$

4 Results and summary

In Fig. 3 we report the cross section curves as functions of $\sqrt{s}$ of the process for the three final states under consideration, $(D\bar{D}, DD^*, D^*D^*)$. This calculation has been made by using the tri- and four-linear couplings quoted above, assuming their validity in the energy range $\sqrt{s} \approx 3.8$–4.5, and computing the tree level diagrams for the process at hand (for a sketch of the diagrams involved, see [1, 2]). The dashed curves define the uncertainties bands obtained by varying Δ and E_ρ, as discussed in Table 1.

We have presented the calculation method of the effective couplings $J/\psi(X)D^{(*)}\bar{D}^{(*)}$, with $X = \rho, \Phi$, within the CQM model. The resulting cross section predictions, together with an estimate of the associated theoretical uncertainties, have been presented as functions of $\sqrt{s}$, showing values of the same order as the cross sections for $J/\psi\pi \to D^{(*)}\bar{D}^{(*)}$. This, given also the higher spin multiplicity of the ρ meson with respect to pions, shows the importance of the ρ contribution to the J/ψ absorption in the hot hadron gas, formed in peripheral heavy-ion collisions at SPS energy, as discussed thoroughly in [1, 2]. Aiming at calculating thermal averages with $T \approx 170$ MeV, we did not discuss in the present paper the introduction of any arbitrary form factors, since the exponential statistical weight acts as a cutoff in the high energy tail.

Acknowledgements. We wish to thank L. Maiani for the stimulating collaboration and encouragement.

Appendix

In this appendix are listed the I_i and L_i integrals occurring in the calculation and their linear combinations A, B, C, D. These integrals have been computed adopting the proper time Schwinger regularization prescription, with cutoff $\mu = 0.3$ GeV (0.5 GeV when a strange quark is present), $\Lambda = 1.25$ GeV. In the following $N_c = 3$. The I_i integrals are

$$\begin{aligned} I_1 &= \mathrm{i}N_c \int \frac{\mathrm{d}^4 l}{(2\pi)^4} \frac{1}{(l^2 - \tilde{m}^2)}\\ &= \frac{N_c}{16\pi^2}\,\tilde{m}^2\,\Gamma\left(-1, \frac{\tilde{m}^2}{\Lambda^2}, \frac{\tilde{m}^2}{\mu^2}\right) \end{aligned} \tag{A.1}$$

$$\begin{aligned} I_3(\delta) &= -\mathrm{i}N_c \int \frac{\mathrm{d}^4 l}{(2\pi)^4} \frac{1}{(l^2 - \tilde{m}^2)\,(v\cdot l + \delta)}\\ &= \frac{N_c}{16\pi^{3/2}} \int_{1/\Lambda^2}^{1/\mu^2} \frac{\mathrm{d}s}{s^{3/2}}\,\mathrm{e}^{-s\left(\tilde{m}^2 - \delta^2\right)}\\ &\quad\times \left(1 + \mathrm{Erf}\left(\delta\sqrt{s}\right)\right)\,, \end{aligned} \tag{A.2}$$

$$\begin{aligned} I_5(\delta, \delta', \omega) &= \mathrm{i}N_c \int \frac{\mathrm{d}^4 l}{(2\pi)^4} \frac{1}{(l^2 - \tilde{m}^2)\,(v\cdot l + \delta)\,(v'\cdot l + \delta')}\\ &= \int_0^1 \mathrm{d}y \frac{1}{1 + 2y^2(1-\omega) + 2y(\omega - 1)}\\ &\quad\times \left[\frac{6}{16\pi^{3/2}} \int_{1/\Lambda^2}^{1/\mu^2} \mathrm{d}s\,\sigma\,\mathrm{e}^{-s\left(\tilde{m}^2 - \sigma^2\right)}\,s^{-1/2}\right.\\ &\quad\left.\times \left(1 + \mathrm{Erf}\left(\sigma\sqrt{s}\right)\right)\right. \end{aligned}$$

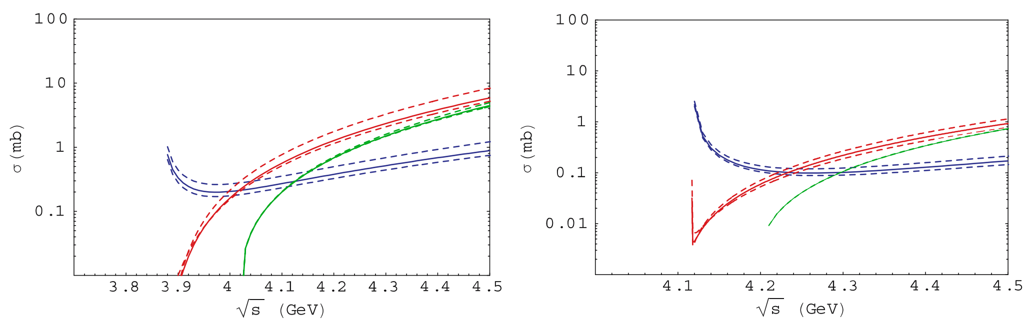

Fig. 3. The cross sections of the processes $J/\psi\,\rho \longrightarrow D^{(*)}\,\bar{D}^{(*)}$ and $J/\psi\,\Phi \longrightarrow D_s^{(*)}\,\bar{D}_s^{(*)}$ *on the left* and *on the right panel*, respectively. The *dashed curves* define the uncertainty bands obtained by varying Δ and E_ρ as discussed above. Some of the reactions, the one initiated by ρ giving $D\bar{D}$ in the final state and those initiated by ϕ giving $D_s\bar{D}_s$ and $D_s^*\bar{D}_s$ (or $D_s\bar{D}_s^*$, a sum of the two is taken), show the typical "exothermic" peak for zero $\rho(\phi)$ momentum. The remaining reactions show the usual threshold behavior (endothermic)

$$+\frac{6}{16\pi^2}\int_{1/\Lambda^2}^{1/\mu^2} \mathrm{d}s\, \mathrm{e}^{-s\sigma^2} s^{-1}\Bigg], \qquad \text{(A.3)}$$

while the L_i integrals are defined as $\frac{\partial}{\partial \tilde{m}^2} I_i$, and they are

$$L_1 = \frac{N_c}{16\pi^2}\left[\Gamma\left(-1, \frac{\tilde{m}^2}{\Lambda^2}, \frac{\tilde{m}^2}{\mu^2}\right) + \tilde{m}^2 \frac{\partial}{\partial \tilde{m}^2}\Gamma\left(-1, \frac{\tilde{m}^2}{\Lambda^2}, \frac{\tilde{m}^2}{\mu^2}\right)\right], \qquad \text{(A.4)}$$

$$L_3(\delta) = \frac{N_c}{16\pi^{3/2}}\int_{1/\Lambda^2}^{1/\mu^2} \mathrm{d}s\, \mathrm{e}^{-s(\tilde{m}^2-\delta^2)}\left(-s^{-1/2}\right) \times \left(1+\mathrm{Erf}\left(\delta\sqrt{s}\right)\right), \qquad \text{(A.5)}$$

$$L_5(\delta,\delta',\omega) = \frac{6}{16\pi^{3/2}}\int_0^1 \mathrm{d}y \frac{1}{1+2y^2(1-\omega)+2y(\omega-1)} \times \int_{1/\Lambda^2}^{1/\mu^2} \mathrm{d}s\, \sigma\, e^{-s(\tilde{m}^2-\sigma^2)}\left(-s^{1/2}\right) \times \left(1+\mathrm{Erf}\left(\sigma\sqrt{s}\right)\right). \qquad \text{(A.6)}$$

In the previous expressions we have defined

$$\sigma \equiv \sigma(\delta,\delta',y,\omega) = \frac{\delta\,(1-y)+\delta'\,y}{\sqrt{1+2\,(\omega-1)\,y+2\,(1-\omega)\,y^2}}, \qquad \text{(A.7)}$$

while $\tilde{m}^2$, δ, δ' and ω are given by (18)–(21). The gamma function and the error function are given respectively by

$$\Gamma(\alpha,x_0,x_1) = \int_{x_0}^{x_1} \mathrm{d}t\, \mathrm{e}^{-t}\, t^{\alpha-1}, \qquad \mathrm{Erf}(z) = \frac{2}{\sqrt{\pi}}\int_0^z \mathrm{d}x\, \mathrm{e}^{-x^2}. \qquad \text{(A.8)}$$

In the last analysis the functions A, B, C, D (which are all functions of x and E_ρ through δ, δ' and ω) are given by

$$A = -\frac{L_3(\delta)+L_3(\delta')+(\delta+\delta')L_5(\delta,\delta',\omega)}{2(1+\omega)}, \qquad \text{(A.9)}$$

$$B = \frac{1}{2(\omega^2-1)}\Big[c\delta' L_3 + \delta L_3' - (\delta L_3 + \delta' L_3')\,\omega + I_5(\omega^2-1) + L_5((\omega^2-1)\tilde{m}^2 + \delta^2 + \delta'^2 - 2\delta\delta'\omega)\Big], \qquad \text{(A.10)}$$

$$C = \frac{1}{2(\omega^2-1)^2}\Big[(2L_1+\delta L_3+\delta' L_3')\,\omega^3 + \big(I_5+\delta L_3'+\delta' L_3 + L_5(\tilde{m}+\delta^2+\delta'^2)\big)\omega^2 2\,(L_1+2\delta L_3+2\delta' L_3'+3\delta\delta' L_5)\,\omega + L_5(2\delta^2+2\delta'^2-\tilde{m}^2) - I_5 + 2\delta L_3' + 2\delta' L_3\Big], \qquad \text{(A.11)}$$

$$D = \frac{1}{2(\omega^2-1)^2}\Big[-\big(L_5\tilde{m}^2+I_5\big)\,\omega^3 + (-2L_1+\delta L_3+\delta' L_3'+4\delta\delta' L_5)\,\omega^2 + \big(L_5(\tilde{m}^2-3\delta^2-3\delta'^2)+I_5-3\delta' L_3-3\delta L_3'\big)\,\omega + 2L_1+2\delta L_3+2\delta' L_3'+2\delta\delta' L_5\Big]. \qquad \text{(A.12)}$$

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