Chiral Lagrangian approach to the $J/\psi$ breakup cross section

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Abstract

We summarize the results of the $SU(4)$ chiral meson Lagrangian approach to the cross section for $J/\psi$ breakup by pion and rho meson impact and suggest a new scheme for the introduction of formfactors for the meson-meson interaction. This scheme respects the fact that on the quark level of description the contact and the meson exchange diagrams are constructed by so-called box and triangle diagrams which contain a different number of vertex functions for the quark-meson coupling. A model calculation for Gaussian vertex functions is presented and the relative importance of contact and meson exchange processes is discussed. We evaluate the dependence of the breakup cross section on the masses of the final D-meson states which can be used for the study of in-medium effects on this quantity.

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I. INTRODUCTION

The $J/\psi$ meson plays a key role in the experimental search for the quark-gluon plasma (QGP) in heavy-ion collision experiments where an anomalous suppression of its production cross section relative to the Drell-Yan continuum as a function of the centrality of the collision has been found by the CERN-NA50 collaboration [1]. An effect like this has been predicted to signal QGP formation [2] as a consequence of the screening of color charges in a plasma in close analogy to the Mott effect (metal-insulator transition) in dense electronic systems [3]. However, a necessary condition to explain $J/\psi$ suppression in the static screening model is that a sufficiently large fraction of $c\bar{c}$ pairs after their creation have to traverse regions of QGP where the temperature (resp. parton density) has to exceed the Mott temperature $T_{Mott}^{J/\psi} \sim 1.2 - 1.3 T_c$ [4,5] for a sufficiently long time interval $\tau > \tau_f$, where $T_c \sim 170$ MeV is the critical phase transition temperature and $\tau_f \sim 0.3$ fm/c is the $J/\psi$ formation time. Within an alternative scenario [6], $J/\psi$ suppression does not require temperatures well above the deconfinement one but can occur already at $T_c$ due to impact collisions by quarks from the thermal medium. An important ingredient for this scenario is the lowering of the reaction threshold for string-flip processes which lead to open-charm meson formation and thus to $J/\psi$ suppression. This process has an analogue in the hadronic world, where e.g. $J/\psi + \pi \rightarrow D^* + \bar{D} + h.c.$ could occur provided the reaction threshold of $\Delta E \sim 640$ MeV can be overcome by pion impact. It has been shown recently [7] that this process and its in-medium modification can play a key role in the understanding of anomalous $J/\psi$ suppression as a deconfinement signal. Since at the deconfinement transition the $D$- mesons enter the continuum of unbound (but strongly correlated) quark-antiquark states (Mott-effect), the relevant threshold for charmonium breakup is lowered and the reaction rate for the process gets critically enhanced. Thus a process which is negligible in the vacuum may give rise to additional (anomalous) $J/\psi$ suppression when conditions of the chiral/deconfinement transition and $D$- meson Mott effect are reached in a heavy-ion collision but the dissociation of the $J/\psi$ itself still needs impact to overcome the threshold which is still present but dramatically reduced.

For this alternative scenario as outlined in [7] to work the $J/\psi$ breakup cross section by pion impact is required and its dependence on the masses of the final state $D$- mesons has to be calculated. Both, nonrelativistic potential models [12,13] and chiral Lagrangian models [9–11] have been employed to determine the cross section in the vacuum. The results of the latter models appear to be strongly dependent on the choice of formfactors for the meson-meson vertices. This is considered as a basic flaw of these approaches which could only be overcome when a more fundamental approach, e.g. from a quark model, can determine these input quantities of the chiral Lagrangian approach.

In the present paper we would like to reduce the uncertainties of the chiral Lagrangian approach by constraining the formfactor from comparison with results of a nonrelativistic approach which makes use of meson wave functions [15]. Finally, we will obtain a result for the off-shell $J/\psi$ breakup cross section which can be compared to the fit formula used in [7]. This quantity is required for the calculation of the in-medium modification of the $J/\psi$ breakup due to the Mott-effect for mesonic states at the deconfinement/chiral restoration transition which has been suggested [4,8] as an explanation of the anomalous $J/\psi$ suppression effect observed in heavy-ion collisions at the CERN-SPS [1].
II. EFFECTIVE CHIRAL LAGRANGIAN

We start from QCD at low energy. The effective chiral Lagrangian for pseudoscalar (Goldstone) mesons can be written as

\[ L_0 = \frac{F_\pi^2}{8} \text{tr} \left[ \partial_\mu U(x) \partial_\mu U^+(x) \right], \]

with \( F_\pi = 132 \text{ MeV} \) being the weak pion decay constant, and \( U(x) = \exp \left[ 2i\varphi(x)/F_\pi \right]. \) Notice that \( U(x) \) transforms in a so-called non-linear representation of the \( SU(N_f)_L \times SU(N_f)_R \) group. The usual multiplet of pseudoscalar mesons is \( \varphi = \varphi^a \lambda_a/\sqrt{2}, \) \( \lambda_a \) are Gell-Mann matrices. To introduce vector and axial-vector mesons we follow the procedure which is connected with the replacement

\[ L_0 \rightarrow L = \frac{F_\pi^2}{8} \text{tr} \left[ D_\mu U D_\mu U^+ \right], \]

given by

\[ D_\mu = \partial_\mu U - igA_\mu^L U + igA_\mu^R. \]

The left- and right-handed spin-1 fields, \( A_\mu^L \) and \( A_\mu^R, \) are combinations of vector and axial-vector meson fields

\[ A_\mu^L = \frac{1}{2} (V_\mu + A_\mu), \]
\[ A_\mu^R = \frac{1}{2} (V_\mu - A_\mu). \]

The coupling of these mesons to pseudoscalars is introduced as a gauge coupling with the gauge coupling constant \( g \) which can be determined from the \( \rho \rightarrow \pi \pi \) decay; \( g_{\rho \pi \pi} = 8.6. \) Therefore, the Lagrangian involving spin-1 and spin-0 mesons takes the form

\[ \mathcal{L}(\varphi, V, A) = \frac{1}{8} F_\pi^2 \text{tr} \left[ D_\mu U (D_\mu U)^+ \right] + \frac{1}{8} F_\pi^2 \text{tr} \left[ M(U + U^+ - 2) \right] \]
\[ - \frac{1}{2} \text{tr} \left[ (F_{\mu \nu}^L)^2 + (F_{\mu \nu}^R)^2 \right] + m_0^2 \text{tr} \left[ (A_{\mu}^L)^2 + (A_{\mu}^R)^2 \right] \]
\[ - i\xi \text{tr} \left[ (D_\mu U)(D_\nu U)^+ F_{\mu \nu}^L + (D_\mu U)^+(D_\nu U)F_{\mu \nu}^R \right] \]
\[ + \gamma \text{tr} \left[ F_{\mu \nu}^L U F_{\mu \nu}^R U^+ \right]. \]

The second term is proportional to the mass matrix \( M \) and describes the “soft” breaking of the chiral \( SU(N_f)_L \times SU(N_f)_R \) symmetry. The corresponding field strength tensors are given by

\[ F_{\mu \nu}^{L,R} = \partial_\mu A_{\nu}^{L,R} - \partial_\nu A_{\mu}^{L,R} - ig \left[ A_{\mu}^{L,R}, A_{\nu}^{L,R} \right]. \]

The third and fourth terms in (5) correspond to the free Lagrangians of the spin-1 particles. At this level of the chiral symmetry all spin-1 mesons have the same “bare” mass, \( m_0. \) The
remaining terms in (5) are so-called non-minimal terms since they contain higher orders in the derivatives as well as the mixed term \((\partial_\mu \varphi A_\mu)\). After the diagonalization of (5) we obtain the Lagrangians with pseudoscalar, vector and axial-vector mesons (see Appendix). The Lagrangian (5) contains many degrees of freedom. We can consider the special case when vector mesons are described as dynamical gauge bosons. It corresponds to the “hidden” chiral symmetry. We choose a gauge where left- and right-handed fields in the Lagrangian will be identical to the vector field \(V_\mu\): \(A^L_\mu = A^R_\mu = V_\mu\) and \(A'_\mu = 0\). This can be achieved by a gauge transformation which conserves the \(SU(N_f)_L \times SU(N_f)_R\) symmetry

\[
A^L_\mu = A^R_\mu = V_\mu , \quad U \rightarrow U_L U U_R^+, \quad A^L_\mu \rightarrow U_L A^L_\mu U_L^+ + \frac{i}{g} U_L \partial_\mu U_R^+, \quad A^L_\mu \rightarrow U_R A^L_\mu U_R^+ + \frac{i}{g} U_R \partial_\mu U_R^+, \tag{7}
\]

with the specific choice \(U_L = U_L^{1/2}\) and \(U_R = U_R^{-1/2}\), so that pseudoscalar mesons are gauge parameters. Now we can rewrite the Lagrangian (5) as a sum of three Lagrangians

\[
\mathcal{L}_0 = \frac{F_\pi^2}{8} \text{tr} \left( D_\mu U D_\mu U^+ \right) , \tag{8}
\]

\[
\mathcal{L}_1 = -\frac{1}{2} \text{tr} \left( (F^L_{\mu\nu})^2 + (F^R_{\mu\nu})^2 \right) + \gamma \text{tr} \left( F^L_{\mu\nu} U F^R_{\mu\nu} U^+ \right) , \tag{9}
\]

\[
\mathcal{L}_2 = m_0^2 \text{tr} \left( (A^L_\mu)^2 + (A^R_\mu)^2 \right) + B \text{tr} (A^L_\mu U A^R_\mu U^+ ) + C \text{tr} (A^L_\mu A^R_\mu + A^R_\mu A^L_\mu) . \tag{10}
\]

Note that we have added two gauge invariant terms to the Lagrangian (5). The second term (with the coefficient B) in (10) plays an important role in the description of the width of the \(\rho \rightarrow \pi\pi\) decay, and the third term (with the coefficient C) maintains the gauge invariance of the \(J/\psi + \pi \rightarrow D^* + \bar{D}\) decay. Applying the substitutions (7) to the Lagrangian (5), we obtain

\[
\mathcal{L}_0 \rightarrow \mathcal{L}'_0 = 0 , \quad \mathcal{L}_1 \rightarrow \mathcal{L}'_1 = (\gamma - 1) \text{tr} \left( F^V_{\mu\nu} F^{\mu\nu V} \right) , \quad \mathcal{L}_2 \rightarrow \mathcal{L}'_2 = \frac{m_0^2}{2} \text{tr} \left( V_{\mu}^2 \right) - i \frac{g V_{\mu\nu\varphi\varphi}}{2} \text{tr} \left( V_{\mu} \left( \varphi_{\mu} \partial_\mu \varphi \right) \right) + \frac{8C}{F_\pi^2} \text{tr} \left( (V_{\mu} \varphi)^2 - V_{\mu}^2 \varphi^2 \right) + \mathcal{L}(\varphi) , \tag{11}
\]

where the vector mass and the vector-pseudoscalar-pseudoscalar coupling are defined by \(m_V^2 = 2(B + 2m_0^2 + 2C)\), \(g_{V\varphi\varphi} = 2(B - 2C + 2m_0^2)/(gF_\pi^2)\).
A. J/ψ breakup cross sections

The above effective Lagrangian allows us to study the following processes for J/ψ breakup by π and ρ mesons

\[
\begin{align*}
J/\psi + \pi &\rightarrow D^* + \bar{D}, \ J/\psi + \pi \rightarrow D + \bar{D}^*, \\
J/\psi + \rho &\rightarrow D + \bar{D}, \ J/\psi + \rho \rightarrow D^* + \bar{D}^*. \quad (12)
\end{align*}
\]

The generic diagrams for these processes are shown in Fig. 1 for the example of the first one.

The full amplitude for the first process \( \psi + \pi \rightarrow D^* + \bar{D}, \ J/\psi + \pi \rightarrow D + \bar{D}^* \), without isospin factors and before summing and averaging over external spins, is given by

\[
M_1 \equiv M_{1 \mu \nu}^\mu \epsilon_1^\mu \epsilon_3^\nu = \left( \sum_{i=a,b,c} M_{1i}^{\mu \nu} \right) \epsilon_1^\mu \epsilon_3^\nu \quad (14)
\]

with

\[
M_{1a}^{\mu \nu} = -g_{\pi DD^*} g_{J/\psi DD^*} (-2 p_2 + p_3)^\nu \left( \frac{1}{u - m_D^2} \right) (p_2 - p_3 + p_4)^\mu,
\]

\[
M_{1b}^{\mu \nu} = g_{\pi DD^*} g_{J/\psi DD^*} (-p_2 - p_4)^\alpha \left( \frac{1}{t - m_{D^*}^2} \right)
\times \left[ g^{\alpha \beta} - \frac{(p_2 - p_4)^\alpha (p_2 - p_4)^\beta}{m_{D^*}^2} \right]
\times \left[ (-p_1 - p_3)^\beta g^{\mu \nu} + (-p_2 + p_1 + p_4)^\nu g^\beta \mu + (p_2 + p_3 - p_4)^\mu g^\beta \nu \right],
\]

\[
M_{1c}^{\mu \nu} = -g_{J/\psi \pi DD^*} g^{\mu \nu}. \quad (15)
\]

Similarly, the full amplitude for the second process \( J/\psi + \rho \rightarrow D + \bar{D} \) is given by

\[
M_2 \equiv M_{2 \mu \nu}^\mu \epsilon_1^\mu \epsilon_2^\nu = \left( \sum_{i=a,b,c} M_{2i}^{\mu \nu} \right) \epsilon_1^\mu \epsilon_2^\nu \quad (16)
\]

with

\[
M_{2a}^{\mu \nu} = -g_{\rho DD} g_{J/\psi DD} (p_2 - 2 p_3)^\mu \left( \frac{1}{u - m_D^2} \right) (p_2 - p_3 + p_4)^\nu,
\]

\[
M_{2b}^{\mu \nu} = -g_{\rho DD} g_{J/\psi DD} (-p_2 + 2 p_4)^\mu \left( \frac{1}{t - m_D^2} \right) (-p_2 - p_3 + p_4)^\nu,
\]

\[
M_{2c}^{\mu \nu} = g_{J/\psi \rho DD} g^{\mu \nu}. \quad (17)
\]

For the third process \( J/\psi + \rho \rightarrow D^* + \bar{D}^* \), the full amplitude is given by
\[ M_3 \equiv M_{3i}^{\mu\nu\lambda\omega} \varepsilon_{1\mu} \varepsilon_{2\nu} \varepsilon_{3\lambda} \varepsilon_{4\omega} = \left( \sum_{i=a,b,c} M_{3i}^{\mu\nu\lambda\omega} \right) \varepsilon_{1\mu} \varepsilon_{2\nu} \varepsilon_{3\lambda} \varepsilon_{4\omega} , \]  

\[ M_{3a}^{\mu\nu\lambda\omega} = g_{\rho D^* D^*} g_{J/\psi D^* D^*} \left[ (-p_2 - p_3)\alpha g^{\mu\lambda} + 2p_3^\mu g^{\alpha\mu} + 2p_3^\rho g^{\alpha\rho} \right] \times \left( \frac{1}{u - m_D^2} \right) \left[ g^{\alpha\beta} - \frac{(p_2 - p_3)\alpha (p_2 - p_3)\beta}{m_D^2} \right] \times \left[ -2p_1^\rho g^{\beta\nu} + (p_1 + p_4)\beta g^{\nu\omega} - 2p_4^\nu g^{\beta\omega} \right] , \]

\[ M_{3b}^{\mu\nu\lambda\omega} = g_{\rho D^* D^*} g_{J/\psi D^* D^*} \left[ -2p_2^\rho g^{\alpha\mu} + (p_2 + p_4)\alpha g^{\mu\omega} - 2p_4^\mu g^{\alpha\omega} \right] \times \left( \frac{1}{t - m_D^2} \right) \left[ g^{\alpha\beta} - \frac{(p_2 - p_4)\alpha (p_2 - p_4)\beta}{m_D^2} \right] \times \left[ (p_1 - p_3)\beta g^{\nu\lambda} + 2p_1^\lambda g^{\beta\nu} + 2p_3^\nu g^{\beta\lambda} \right] , \]

\[ M_{3c}^{\mu\nu\lambda\omega} = g_{J/\psi p D^* D^*} (g^{\mu\lambda} g^{\nu\omega} + g^{\mu\omega} g^{\nu\lambda} - 2g^{\mu\nu} g^{\lambda\omega}) . \]  

In the above, \( p_j \) denotes the momentum of particle \( j \). We choose the convention that particle 1 and 2 represent initial-state mesons while particles 3 and 4 represent final-state mesons on the left and right sides of the diagrams shown in Fig. 1, respectively. The indices \( \mu, \nu, \lambda \) and \( \omega \) denote the polarization components of external particles while the indices \( \alpha \) and \( \beta \) denote those of the exchanged mesons.

After averaging (summing) over initial (final) spins and including isospin factors, the cross sections for the three processes are given by

\[ \frac{d\sigma_1}{dt} = \frac{1}{288 \pi s p_{\text{c.m.}}^2} M_1^{\mu\nu} M_1^{\nu'\mu'} \left( g^{\mu\nu} - \frac{p_1^\mu p_1^{\nu'}}{m_1^2} \right) \left( g^{\nu\nu'} - \frac{p_2^\nu p_2^{\nu'}}{m_2^2} \right) , \]

\[ \frac{d\sigma_2}{dt} = \frac{1}{288 \pi s p_{\text{c.m.}}^2} M_2^{\mu\nu} M_2^{\nu'\mu'} \left( g^{\mu\nu} - \frac{p_1^\mu p_1^{\nu'}}{m_1^2} \right) \left( g^{\nu\nu'} - \frac{p_2^\nu p_2^{\nu'}}{m_2^2} \right) , \]

\[ \frac{d\sigma_3}{dt} = \frac{1}{288 \pi s p_{\text{c.m.}}^2} M_3^{\mu\nu\lambda\omega} M_3^{\nu'\mu'\lambda'\omega'} \left( g^{\mu\nu} - \frac{p_1^\mu p_1^{\nu'}}{m_1^2} \right) \left( g^{\nu\nu'} - \frac{p_2^\nu p_2^{\nu'}}{m_2^2} \right) \times \left( g^{\lambda\lambda'} - \frac{p_3^\lambda p_3^{\lambda'}}{m_3^2} \right) \left( g^{\omega\omega'} - \frac{p_4^\omega p_4^{\omega'}}{m_4^2} \right) , \]

with \( s = (p_1 + p_2)^2 \), and

\[ p_{\text{c.m.}}^2 = \frac{[s - (m_1 + m_2)^2][s - (m_1 - m_2)^2]}{4s} , \]

is the squared momentum of initial-state mesons in the center-of-momentum (c.m.) frame. The definition of \( p_{\text{f.c.m.}} \) for the final-state mesons is analogous with the replacement \((m_1, m_2) \rightarrow (m_3, m_4)\).
III. HADRONIC FORMFACTORS

The chiral Lagrangian approach for $J/\psi$ breakup by light meson impact makes the assumption that mesons and meson-meson interaction vertices are pointlike (four-momentum independent) objects. This neglect of the finite extension of mesons as quark-antiquark bound states has dramatic consequences: it leads to a monotonically rising behaviour of the cross sections for the corresponding processes, see the dashed lines in Fig. 2. This result, however, cannot be correct away from the reaction threshold where the tails of the mesonic wave functions determine the high-energy behaviour of the quark exchange (in the nonrelativistic formulation of [12,15]) or quark loop (in the relativistic formulation [16]) diagrams describing the microscopic processes underlying the $J/\psi$ breakup by meson impact. As long as the mesonic wave functions describe quark-antiquark bound states which have a finite extension in coordinate- and momentum space, the $J/\psi$ breakup cross section is expected to be decreasing above the reaction threshold and asymptotically small at high c.m. energies. This result of the quark model approaches to meson-meson interactions [12,15,16] can be mimicked within chiral meson Lagrangian approaches by the use of formfactors at the interaction vertices [10,11].

A. Global formfactor ansatz

We will follow here the definitions of Ref. [10], where the formfactor of all the four-point vertices of Fig.1, i.e. that of the box diagram (I) as well as that of the meson exchange diagrams (II, III) is taken as a product of the triangle diagram formfactors

$$F^i_4(q^2) = \left[F_3(q^2)\right]^2, \quad i = I, II, III,$$

with the squared three-momentum $q^2$ given by the average value of the squared three-momentum transfers in the $t$ and $u$ channels

$$q^2 = \frac{1}{2} \left[ (p_1 - p_3)^2 + (p_4 - p_2)^2 \right]_{c.m.} = p_{i,c.m.}^2 + p_{f,c.m.}^2.$$

For the triangle diagrams, we use formfactors with a momentum dependence in the monopole form ($M$)

$$F^M_3(q^2) = \frac{\Lambda^2}{\Lambda^2 + q^2},$$

and in the Gaussian ($G$) form

$$F^G_3(q^2) = \exp\left(-q^2/\Lambda^2\right).$$
B. Meson formfactor ansatz

In order to take into account the quark substructure of meson-meson vertices we will suggest here a simple ansatz which respects the different size of the interacting mesons and the different quark diagram representation of contact and meson exchange interactions in terms of quark box and quark triangle diagrams. The triangle diagram is of third order in the wave functions so that the meson exchange diagrams are suppressed at large momentum transfer by six wave functions, the box diagram appears already at fourth order thus being less suppressed than suggested by the ansatz \[^{[24]}\] of Ref. \[^{[10]}\].

For the contact term (I) we use the replacement \(g_{J/ψπD^∗D} \rightarrow g_{J/ψπD^∗D} \times F_I(s)\) where the formfactor has the following form

\[
F_I(s) = \exp \left\{ -\frac{1}{4s} \left[ (s - (m_1 + m_2)^2) (s - (m_1 - m_2)^2) \left( \frac{1}{Λ_1^2} + \frac{1}{Λ_2^2} \right) \\
+ (s - (m_3 + m_4)^2) (s - (m_3 - m_4)^2) \left( \frac{1}{Λ_3^2} + \frac{1}{Λ_4^2} \right) \right] \right\} \tag{28}
\]

The second exchange term (II) can be written using the substitution \(g_{J/ψD^∗D^∗} \times g_{D^∗Dπ} \rightarrow g_{J/ψD^∗D^∗} \times g_{D^∗Dπ} \times F_{II}(s,t)\) with

\[
F_{II}(s,t) = \exp \left\{ -\frac{1}{4s} \left[ (s - (m_1 + m_2)^2) (s - (m_1 - m_2)^2) \left( \frac{1}{Λ_1^2} + \frac{1}{Λ_2^2} \right) \\
+ (s - (m_3 + m_4)^2) (s - (m_3 - m_4)^2) \left( \frac{1}{Λ_3^2} + \frac{1}{Λ_4^2} \right) \\
+ \frac{2}{Λ_1^2} \left[ \left( m_1^2 - m_2^2 \right) - \left( m_3^2 - m_4^2 \right) \right] - 4st \right] \right\} \tag{29}
\]

Analogously, the exchange term (III) is obtained by \(g_{J/ψDD} \times g_{D^∗Dπ} \rightarrow g_{J/ψDD} \times g_{D^∗Dπ} \times F_{III}(s,u)\) with the formfactor

\[
F_{III}(s,u) = \exp \left\{ -\frac{1}{4s} \left[ (s - (m_1 + m_2)^2) (s - (m_1 - m_2)^2) \left( \frac{1}{Λ_1^2} + \frac{1}{Λ_2^2} \right) \\
+ (s - (m_3 + m_4)^2) (s - (m_3 - m_4)^2) \left( \frac{1}{Λ_3^2} + \frac{1}{Λ_4^2} \right) \\
+ \frac{2}{Λ_1^2} \left[ \left( m_1^2 - m_2^2 \right) + \left( m_3^2 - m_4^2 \right) \right] - 4su \right] \right\} \tag{30}
\]

Formfactors depend on the parameters which we fix from the physical observables (decay widths \(ρ \rightarrow ππ, D^∗ \rightarrow Dπ\) and the vector dominance model \[^{[10]}\][^{[11]}\]. We use for the coupling constants the values \(g_{D^∗Dπ} = g_{DDπ} = g_{D^∗D^∗π} = 4.4, g_{J/ψDD} = g_{J/ψD^∗D^∗} = 7.7, g_{J/ψπD^∗D} = g_{J/ψπD^∗π} = 33.9\). For the range parameters \(Λ_i\) of the quark-antiquark-meson vertices we suggest to use the meson masses \(m_i\), see Table \[^{[15]}\]. The results are depicted in Fig. \[^{[2]}\]. In the last Section, we discuss the results and their possible implications for phenomenological applications.
IV. RESULTS AND DISCUSSION

The $J/\psi$ breakup cross section by $\pi$ and $\rho$ meson impact has been formulated within a chiral $U(4)$ Lagrangian approach. Numerical results have been obtained for the pion impact processes with the result that the D-meson exchange in the t-channel is the dominant subprocess contributing to the $J/\psi$ breakup. The use of formfactors at the meson-meson vertices is mandatory since otherwise the high-energy asymptotics of the processes with hadronic final states will be overestimated, see Fig. 2. From a comparison of the monopole and Gaussian functions in the global formfactor ansatz, we observe a big difference in the corresponding cross sections above the threshold. Generally we would prefer the use of gaussian functions which are motivated by strong (confining) quark-antiquark interactions within the mesons. The new meson formfactor scheme reduces the peak value of the $J/\psi$ breakup cross section relative to the global scheme by 50%. The net result for the considered pion impact subprocess is in good correspondence to the one from the nonrelativistic quark exchange model.

Finally, we want to present an exploratory study of the influence of a variation of the final state D-meson masses on the effective $J/\psi$ breakup cross section. Our motivation for considering mesonic states to be off their mass-shell is their compositeness which can become apparent in a high-temperature (and density) environment at the deconfinement/chiral restoration transition, when these states change their character qualitatively being resonant quark-antiquark scattering states in the quark plasma rather than on-shell mesonic bound states.

The consequence of this Mott-transition from bound to resonant states for the $J/\psi$ breakup has been explored by Burau et al. [7,8], using a fit formula for the D-mass dependence of the breakup cross section which shows a strong enhancement when the process becomes subthreshold. This behaviour is qualitatively approved within the present chiral $U(4)$ Lagrangian + formfactor model although the subthreshold enhancement is more moderate, see Fig. 2. A more consistent description should include a quark model derivation of the formfactors for the meson-meson vertices and their possible medium dependence. Such an investigation is in progress.

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APPENDIX A: CHIRAL LANGRANGIANS

In this Appendix we write the explicit form of the chiral Langrangians. We use these Langrangians for the calculations of matrix elements of the $J/\psi$ breakup cross section by $\pi$ and $\rho$ mesons
\[ \mathcal{L}^{(2)}(\varphi, V, A) = \frac{1}{2} \text{tr}(\partial_\mu \varphi)^2 - \frac{1}{2} \text{tr}(M \varphi^2) - \frac{1}{4} \text{tr}(F_{\mu\nu}^V)^2 + \frac{1}{2} m_\nu^2 \text{tr}(V_\mu)^2 - \frac{1}{4} \text{tr}(F_{\mu\nu}^A)^2 + \frac{1}{2} m_\nu^2 \text{tr}(A_\mu)^2 , \]  
(A1)

\[ \mathcal{L}(\varphi^4) = -\frac{2}{3F_\pi^2} \text{tr}(\partial_\mu \varphi \partial_\mu (\varphi^3)) + \frac{1}{2F_\pi^2} Z^2 \text{tr}(\partial_\mu (\varphi^2) \partial_\mu (\varphi^2)) + \frac{g^2}{4m_\nu^2} \text{tr}(\partial_\mu \varphi (\varphi^2, \partial_\mu \varphi)) + \frac{1}{6F_\pi^2} \text{tr}(M \varphi^4) + \left( \frac{1}{8} g^2 (1 - \gamma)^2 \alpha^4 - \frac{2g}{F_\pi^2} Z^4 \alpha^2 \sqrt{1 - \gamma} \right) \text{tr}(\partial_\mu \varphi \partial_\nu \varphi [\partial_\mu \varphi, \partial_\nu \varphi]) , \]  
(A2)

\[ \mathcal{L}(VV\varphi \varphi) = -\frac{g^2}{4Z^2} \text{tr} \left( (V_\mu \varphi)^2 - V_\mu^2 \varphi^2 \right) - \frac{1}{F_\pi^2} \frac{\gamma}{1 - \gamma} \text{tr} \left( (\varphi^2 (F_{\mu\nu}^V)^2 - (F_{\mu\nu}^V \varphi)^2 \right) + \frac{1}{16} g^2 \alpha^2 (1 + \gamma) \text{tr} \left( [\partial_\mu \varphi, V_\nu] + [V_\mu, \partial_\nu \varphi] \right)^2 + \frac{1}{8} g^2 \alpha^2 (1 - \gamma) \text{tr} \left( [V_\mu, V_\nu] [\partial_\mu \varphi, \partial_\nu \varphi] \right) + \frac{g\alpha}{2F_\pi} \frac{\gamma}{\sqrt{1 - \gamma}} \text{tr} \left( \varphi [F_{\mu\nu}^V, ([\partial_\mu \varphi, V_\nu] + [V_\mu, \partial_\nu \varphi])] \right) - \frac{2g\xi}{F_\pi} \frac{Z^4}{\sqrt{1 - \gamma}} \text{tr} (\partial_\mu \varphi \partial_\nu \varphi [V_\mu, V_\nu]) + \frac{2g\xi}{F_\pi} \frac{Z^2}{\sqrt{1 - \gamma}} \text{tr} \left( (\partial_\mu \varphi [V_\nu] + [\varphi, V_\mu] \partial_\nu \varphi) F_{\mu\nu}^V \right) , \]  
(A3)

\[ \mathcal{L}(A, V, \varphi) = -i \frac{g^2 F_\pi}{4Z^2} \sqrt{\frac{1 - \gamma}{1 + \gamma}} \text{tr} (V_\mu [A_\mu, \varphi]) + i \frac{\gamma}{F_\pi} \sqrt{\frac{1 - \gamma}{1 - \gamma}} \text{tr} \left( \varphi [F_{\mu\nu}^V, F_{\mu\nu}^A] \right) + i \frac{g^2 F_\pi}{4m_\nu^2 Z^2} (1 - \delta) \sqrt{\frac{1 - \gamma}{1 + \gamma}} \text{tr} \left( F_{\mu\nu}^V [A_\mu, \partial_\nu \varphi] \right) + i \frac{g^2 F_\pi}{4m_\nu^2} \sqrt{\frac{1 + \gamma}{1 - \gamma}} \text{tr} \left( F_{\mu\nu}^A [V_\mu, \partial_\nu \varphi] \right) , \]  
(A4)

\[ \mathcal{L}(V\varphi \varphi) = -i \frac{g}{2} \text{tr} \left( V_\mu \left( \varphi \stackrel{\leftrightarrow}{\partial_\mu} \varphi \right) \right) + i \frac{g\delta}{2m_\nu^2} \text{tr} \left( F_{\mu\nu}^V \partial_\mu \varphi \partial_\nu \varphi \right) , \]  
(A5)

\[ \mathcal{L}(V^4) = \frac{1}{16} \frac{g^2}{1 - \gamma} \text{tr} \left( [V_\mu, V_\nu]^2 \right) , \]  
(A6)
\[ \mathcal{L}(V^3) = i \frac{g}{4} \text{tr} \left( F_{\mu\nu} [V_\mu, V_\nu] \right) , \] 
\[ \text{(A7)} \]

where
\[ \delta = 1 - Z^2 - \frac{2Z^4}{1 - Z^2} \frac{\xi g}{\sqrt{1 - \gamma}} , \]
and \[ F^V_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu, \quad F^A_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu, \quad (\varphi \stackrel{\leftrightarrow}{\partial_\mu} \varphi) = \varphi \partial_\mu \varphi - \partial_\mu \varphi \varphi , \]
which have the standard definition for commutators and anticommutators \[ [17]. \]
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| state $i$ | $J/\psi$ | $D^*$ | $D$ | $\varrho$ | $\pi$ |
|-----------|-----------|-------|-----|-----------|-------|
| $m_i[\text{GeV}]$ | 3.1       | 2.01  | 1.87| 0.77      | 0.14  |
| $\Lambda_i[\text{GeV}]$ | 3.1       | 2.0   | 1.9 | 0.8       | 0.6   |

**TABLE I.** Meson masses and range parameters corresponding to the quark-antiquark-meson vertices.
FIG. 1. Diagrams for $J/\psi$ breakup by pion impact: $J/\psi + \pi \rightarrow D^* + \bar{D}$; I - contact term, II+III - D-meson exchange processes.
FIG. 2. Upper right panel: total cross section for $J/\psi$ break-up by pion impact without formfactor (black dashed line), with monopole type formfactor (black dotted-dashed-dotted line), with Gaussian formfactor (black dotted-dash line) and with “meson” formfactor (red solid line) as a function of the c.m. energy of initial-state mesons. The partial contributions from the diagrams I, II, and III of Fig. [I] are shown in the other panels.
FIG. 3. Total $J/\psi + \pi \rightarrow \bar{D} + D^*$ cross section with our meson formfactor (black solid line) in comparison with the cross sections of all sub-processes (contact term, exchange terms and corresponding interference terms)
FIG. 4. Total cross section for $J/\psi$ break-up by $\rho$-meson impact due to the exothermic process $J/\psi + \rho \rightarrow \bar{D} + D$ without formfactor (black dashed line) and with meson formfactor (red solid line) as a function of the c.m. energy of initial-state mesons.
FIG. 5. The same like in Fig. 4, but due to the endothermic process $J/\psi + \rho \rightarrow \bar{D}^* + D^*$ without formfactor (black dashed line) and with meson formfactor (red solid line) as a function of the c.m. energy of initial-state mesons. A representation with linear scaling is shown in the small panel.
FIG. 6. Total $J/\psi$ break-up cross section in the chiral Lagrangian model with mesonic form-factor when the final state masses $M_{D_1} = M_{D_2} = M_D$ are varied (“off-shell”) in the way that the reaction threshold for the process $J/\psi + \pi \rightarrow \bar{D} + D^*$ is decreasing (endothermic; black dotted-dashed line) and finally vanishing (exothermic; black dashed line). The red solid line shows the “on-shell” result (the same like in Fig. 3).