In relativistic heavy ion collisions $J/\psi$ suppression has been recognized as an important tool to identify the possible phase transition to quark-gluon plasma (QGP). Matsui and Satz [1] predicted that in presence of quark-gluon plasma, binding of a $c\bar{c}$ pair into a $J/\psi$ meson will be hindered, leading to the so called $J/\psi$ suppression in heavy ion collisions. Over the years several experiments measured the $J/\psi$ yield in heavy ion collisions (for a review of data and interpretations see Ref. [2, 3]). In brief, experimental data do show suppression [4]. However, this could be attributed to more conventional $J/\psi$ absorption by comovers, not present in pA collisions [3, 4, 7]. In order to confirm that the suppression of $J/\psi$ comes from the presence of the QGP, it is necessary to understand better the $J/\psi$ dissociation mechanism by collision with comoving hadrons.

Since there is no empirical information on $J/\psi$ absorption cross sections by hadrons, theoretical models are needed to estimate their values. In general, different models apply to different energy regimes and one of the first estimates of the charmonium-hadron cross section uses short distance QCD [8, 9, 10]. However, the method is inapplicable at low energies, which is the regime of greatest interest for $J/\psi$ collision with comoving hadrons. Besides, even in the high energy regime, nonperturbative effects may be important [11] and can increase significantly the value of the cross section. At the low energy regime one can use quark-interchange models [12] or meson exchange models [13, 14]. The results of the calculations for the charmonium-pion cross section based on these two approaches can differ by two orders of magnitude in the relevant energy range. Moreover, the rapid increase of the cross section near the threshold, obtained with these two models, is probably overestimated since these models do not respect chiral symmetry, as showed in [14]. There is also a calculation of the $J/\psi - \pi$ cross section [15] based on the QCD sum rules (QCDSR) technique [16, 17], which is also valid at the low energy regime. The result for the cross section in ref. [15] is in between the results in the quark-interchange models and meson exchange models.

In this work we improve the calculation done in ref. [15] by considering sum rules based on a three-point function with a pion. We work up to twist-4, which allows us to study the convergence of the OPE expansion. Since the method of the QCDSR uses QCD explicitly, we believe that our work will improve the understanding of this important topic.

In the QCDSR approach, the short range perturbative QCD is extended by an OPE expansion of the correlator, giving a series in inverse powers of the squared momentum with Wilson coefficients. The convergence at low momentum is improved by using a Borel transform. The coefficients involve 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, giving sum rules from which a hadronic quantity can be estimated.

Let us start with the vacuum-pion correlation function for the process $J/\psi \pi \rightarrow D D$:

$$
\Pi_{\mu} = \int d^{4}x \, d^{4}y \, e^{-ip_{2} \cdot y} \, e^{ip_{3} \cdot x} \times \langle 0 | T \{ j_{D}(x) j_{D}(0) j_{\mu}^{\pi}(y) \} | \pi(p_{1}) \rangle \, ,
$$

(1)

with the currents given by $j_{\pi}^{\mu} = \bar{c} \gamma_{\mu} c$, $j_{D} = \bar{c} \gamma_{\mu} D_{\mu}$, $p_{1}$, $p_{2}$, $p_{3}$ and $p_{4}$ are the four-momenta of the mesons $\pi$, $J/\psi$, $D$ and $\bar{D}$ respectively. The advantage of this approach as compared with the 4-point calculation in ref. [15], is that we can consider more terms in the OPE expansion of the correlation function in Eq. (1) and, therefore, we get a much richer sum rule.

Following ref. [16], we can rewrite Eq. (1) as:

$$
\Pi_{\mu} = \int \frac{d^{4}k}{(2\pi)^{4}} \, Tr[S(p_{3} - k) \gamma_{\mu} \times S(p_{3} - p_{2} - k) \gamma_{5} D_{ab}(k, p_{1}) \gamma_{5}] \, ,
$$

(2)

where $S(p)$ is the free c-quark propagator, and $D_{ab}(k, p)$ denotes the quark-antiquark component with a pion, which can be separated into three pieces depending on the Dirac matrices involved [16]:

$$
D_{ab}(k, p) = \delta_{ab} \left[ \gamma_{5} A + \gamma_{\alpha} \gamma_{5} B^{\alpha} + \gamma_{5} \sigma_{\alpha\beta} C^{\alpha\beta} \right] \, .
$$

(3)

The three invariant functions of $(k, p)$: $A$, $B^{\alpha}$ and $C^{\alpha\beta}$, are defined by the Fourier transform of the...
vacuum-pion matrix elements: \( \langle 0|\bar{d}(x)i\gamma_5u(0)|\pi(p_1)\rangle \), \( \langle 0|\bar{d}(x)\gamma^\alpha\gamma_5u(0)|\pi(p_1)\rangle \), and \( \langle 0|\bar{d}(x)\sigma^{\alpha\beta}\gamma_5u(0)|\pi(p_1)\rangle \) respectively.

Using PCAC and working at the order \( \mathcal{O}(p_\mu p_\nu) \) we get up to twist-4

\[
A(k, p) = \frac{(2\pi)^4}{12} \frac{i\langle \bar{q}q \rangle}{f_\pi} \left[ -2 + ip_{\alpha_1} \frac{\partial}{\partial q_{\alpha_1}} + \frac{1}{2} \left( -\frac{m_0^2}{4} \right) \right. \\
\left. \times g_{\alpha_1\alpha_2} + \frac{2}{3} p_{\alpha_1} p_{\alpha_2} \frac{\partial}{\partial q_{\alpha_1}} \frac{\partial}{\partial q_{\alpha_2}} \right] \delta^{(4)}(k), \\
B_\alpha(k, p) = \frac{(2\pi)^4 f_\pi}{12} \left[ ip_{\alpha} + \frac{1}{2} p_{\alpha} p_{\alpha_1} \frac{\partial}{\partial q_{\alpha_1}} + \frac{i\delta^2}{36} \right. \\
\left. \times 5p_\alpha g_{\alpha_1\alpha_2} - 2p_{\alpha_2} g_{\alpha_1} \right] \delta^{(4)}(k), \\
C_{\alpha\beta}(k, p) = \frac{(2\pi)^4}{24} \frac{i\langle \bar{q}q \rangle}{3f_\pi} \left[ p_\alpha g_{\alpha_1\alpha} - p_\beta g_{\alpha_1\alpha} \right] \\
\times \left[ i \frac{\partial}{\partial q_{\alpha_1}} - \frac{p_{\alpha_2}}{2} \frac{\partial}{\partial q_{\alpha_2}} \right] \delta^{(4)}(k),
\]

where \( m_0^2 \) and \( \delta^2 \) are defined by \( \langle \bar{q}D^2 q \rangle = m_0^2 \langle \bar{q}q \rangle / 2 \), \( \langle 0|g_\mu g_\nu\bar{u}(0)u(p)\rangle = i\delta^2 f_\pi p^\mu p^\nu \), with \( g_{\alpha\beta} = \epsilon_{\alpha\beta\gamma\delta}G^{\gamma\delta}/2 \) and \( G_{\alpha\beta} = i^4 G_{\alpha\beta} \).

The phenomenological side of the correlation function, \( \Pi_\mu \), is obtained by the consideration of \( J/\psi \), \( \pi \), and \( D \) states contribution to the matrix element in Eq. (1). The hadronic amplitude is defined by the matrix element:

\[
i\mathcal{M} = \langle \psi(p_2, \mu)| D(\to) D(\to) \pi(p_1)\rangle = i\mathcal{M}_\mu(p_1, p_2, p_3, p_4) e_{\mu}^2.
\]

The phenomenological side of the sum rule can be written as (for the part of the hadronic amplitude that will contribute to the cross section)

\[
\Pi^{\text{phen}} = -\frac{m_\psi f_\psi(m_\psi^2 f_{D/m_\psi^2})^2 M_\mu}{(p_2^2 - m_\psi^2)(p_3^2 - m_\psi^2)(p_4^2 - m_D^2)} + h. \ r. ,
\]

where \( h. \ r. \) means higher resonances. The hadronic amplitude can be parametrized as:

\[
\mathcal{M}_\mu = \Lambda \epsilon_{\mu\alpha\beta\gamma} p_1^\alpha p_3^\beta p_4^\gamma,
\]

where \( \Lambda \) is the parameter that we will evaluate from the sum rules.

Inserting the results in Eqs. (1) and (7) into Eqs. (2) and (3) we can write a sum rule for the invariant structure appearing in Eq. (10). To improve the matching between the phenomenological and theoretical sides we follow the usual procedure and make a single Borel transformation to all the external momenta taken to be equal: \( -p_2^2 = -p_3^2 = -p_4^2 = D^2 \to M^2 \). We get, in the approximation \( p_1 < p_2, p_3, p_4 \):

\[
\frac{\lambda + AM^2 + BM^4}{m_\psi^2 - m_D^2} \left[ e^{-m_\psi^2/M^2} - e^{-m_D^2/M^2} - e^{-m_\psi^2/M^2} \right] \\
\frac{m_\psi^2 - m_D^2}{m_\psi^4 f_\psi f_\pi} \left[ f_\pi \left( \frac{2m_\psi \langle \bar{q}q \rangle}{3f_\pi M^2} \right) \\
- \frac{f_\pi \delta^2}{18M^2} \left( 17 + \frac{5m_\psi^2}{M^2} \right) \right],
\]

where we have transferred to the theoretical side the couplings of the currents with the mesons. The problem of doing a single Borel transformation is the fact that terms associated with the pole-continuum transitions are not suppressed. In the present case we have two kinds of these transitions: double pole-continuum and single pole-continuum. In the limit of similar meson masses it is easy to show that the Borel behavior of the three-pole, double pole-continuum and single pole-continuum contributions are \( m_\psi^2 M^2 / M^4 \), \( m_\psi^2 M^2 / M^2 \), and \( m_\psi^2 M^2 / M^2 \) respectively. Therefore, we can single out the three-pole contribution from the others by introducing two parameters, \( A \) and \( B \), in the phenomenological side of the sum rule, which will account for the double pole-continuum and single pole-continuum contributions respectively.

The parameter values used in all calculations are \( m_\psi = 1.37 \text{ GeV}, m_\pi = 140 \text{ MeV}, m_D = 1.87 \text{ GeV}, m_{D^*} = 2.01 \text{ GeV}, m_\psi = 3.097 \text{ GeV}, f_\pi = 131.5 \text{ MeV}, \)
f_{\psi} = 270 \text{ MeV}, \quad f_D = 170 \text{ MeV}, \quad f_{D^*} = 240 \text{ MeV}, \quad \langle \bar{q} q \rangle = -(0.23)^3 \text{ GeV}^3, \quad m_0 = 0.8 \text{ GeV}^2, \quad \delta^2 = 0.2 \text{ GeV}^2, \quad f_{3\pi} = 0.0035 \text{ GeV}^2. \#10

In Fig. 1 we show the QCD sum rule results for \( \Lambda + AM^2 + BM^4 \) as a function of \( M^2 \). The circles, squares and diamonds give the twist-2, 3 and 4 contributions respectively. The triangles give the final QCDSR results. We see that the twist-3 and 4 contributions are small as compared with the twist-2 contribution showing a “convergence” of the OPE expansion. The triangles follow almost a straight line in the Borel region \( 6 \leq M^2 \leq 16 \text{ GeV}^2 \), indicating that the single pole-continuum transitions contribution is small. The value of the amplitude \( \Lambda \) is obtained by the extrapolation of the fit to \( M^2 = 0 \). \#20, \#21, \#22. Fitting the QCD sum rule results to a quadratic form we get \( \Lambda \approx 11.4 \text{ GeV}^{-3} \). As expected, in our approach \( \Lambda \) is just a number and all dependence of \( M_\mu \) on particle momenta is contained in the Dirac structure. This is a consequence of our low energy approximation.

Instead of using the experimental values for the meson decay constants, it is also possible to use the respective sum rules, as done in \#13. The behavior of the results does not change significantly, leading only to a change in the value of the amplitude. Using the respective sum rules for the meson decay constants we get \( \Lambda \approx 14.9 \text{ GeV}^{-3} \). We will use these two procedures to estimate the errors in our calculation. Our results agrees completely with the value obtained in \#13.

The calculation of the sum rules for the processes \( J/\psi \pi \rightarrow \bar{D} \bar{D}^* \) and \( J/\psi \pi \rightarrow \bar{D}^* \bar{D}^* \) can be done in a similar way. One has only to change the currents in Eq. \#11 by the appropriate ones. The hadronic amplitudes for these two processes can be written in terms of many different structures. In terms of the structures that will contribute to the cross section we can write:

- for the process \( J/\psi \pi \rightarrow \bar{D} \bar{D}^* \):

\[
M_{\mu\nu} = \Lambda_1^{D\bar{D}^*} p_{1\mu} p_{1\nu} + \Lambda_2^{D\bar{D}^*} p_{1\mu} p_{2\nu} + \Lambda_3^{D\bar{D}^*} p_{1\mu} p_{3\nu} + \Lambda_4^{D\bar{D}^*} g_{\mu\nu} + \Lambda_5^{D\bar{D}^*} p_{2\mu} p_{3\nu},
\]

- for the process \( J/\psi \pi \rightarrow \bar{D}^* \bar{D}^* \):

\[
M_{\mu\nu} = \Lambda_1^{D^*\bar{D}^*} H_{\mu\nu} + \Lambda_2^{D^*\bar{D}^*} J_{\mu\nu} + \Lambda_3^{D^*\bar{D}^*} \epsilon_{\rho\sigma\beta\gamma} p_{1\rho} p_{2\sigma} p_{3\beta} + \Lambda_4^{D^*\bar{D}^*} \epsilon_{\rho\sigma\beta\gamma} p_{3\rho} p_{1\sigma} p_{2\beta} + \Lambda_5^{D^*\bar{D}^*} \epsilon_{\rho\sigma\beta\gamma} p_{2\rho} p_{1\sigma} p_{3\beta} + \Lambda_6^{D^*\bar{D}^*} \epsilon_{\rho\sigma\beta\gamma} p_{1\rho} p_{3\sigma} p_{2\beta} + \Lambda_7^{D^*\bar{D}^*} \epsilon_{\rho\sigma\beta\gamma} p_{2\rho} p_{3\sigma} p_{1\beta} + \Lambda_8^{D^*\bar{D}^*} \epsilon_{\rho\sigma\beta\gamma} p_{1\rho} p_{2\sigma} p_{3\beta} + \Lambda_9^{D^*\bar{D}^*} \epsilon_{\rho\sigma\beta\gamma} p_{1\rho} p_{3\sigma} p_{2\beta} + \Lambda_{10}^{D^*\bar{D}^*} \epsilon_{\rho\sigma\beta\gamma} p_{2\rho} p_{3\sigma} p_{1\beta} + \Lambda_{11}^{D^*\bar{D}^*} \epsilon_{\rho\sigma\beta\gamma} p_{3\rho} p_{1\sigma} p_{2\beta} + \Lambda_{12}^{D^*\bar{D}^*} \epsilon_{\rho\sigma\beta\gamma} p_{3\rho} p_{2\sigma} p_{1\beta} + \Lambda_{13}^{D^*\bar{D}^*} \epsilon_{\rho\sigma\beta\gamma} p_{3\rho} p_{1\sigma} p_{2\beta} + \Lambda_{14}^{D^*\bar{D}^*} \epsilon_{\rho\sigma\beta\gamma} p_{1\rho} p_{3\sigma} p_{2\beta},
\]

where

- \( H_{\mu\nu} = \left( \epsilon_{\rho\sigma\beta\gamma} g_{\mu\nu} - \epsilon_{\rho\sigma\beta\gamma} \epsilon_{\mu\rho\beta\gamma} p_{1\nu} p_{2\rho} p_{3\beta} \right) + \epsilon_{\rho\sigma\beta\gamma} \epsilon_{\mu\rho\beta\gamma} p_{1\nu} p_{2\rho} p_{3\beta} + \epsilon_{\rho\sigma\beta\gamma} \epsilon_{\mu\rho\beta\gamma} p_{1\nu} p_{2\rho} p_{3\beta} \) and
- \( J_{\mu\nu} = \left( \epsilon_{\rho\sigma\beta\gamma} \epsilon_{\mu\rho\beta\gamma} + \epsilon_{\rho\sigma\beta\gamma} \epsilon_{\mu\rho\beta\gamma} + \epsilon_{\rho\sigma\beta\gamma} \epsilon_{\mu\rho\beta\gamma} \right) p_{1\nu} p_{2\rho} p_{3\beta} + \epsilon_{\rho\sigma\beta\gamma} \epsilon_{\mu\rho\beta\gamma} p_{1\nu} p_{2\rho} p_{3\beta}. \) In principle all the independent structures appearing in \( H_{\mu\nu} \) and \( J_{\mu\nu} \) would have independent parameters \( \Lambda_i \). However, since in our approach we get exactly the same sum rules for all of them, we decided to group them with the same parameters.

The expressions for all 20 sum rules will be given elsewhere \#23. At this point it is important to stress that the sum rule for the process \( J/\psi \pi \rightarrow \bar{D} \bar{D} \) is not particular, in general all the other sum rules are similar and contain twist-2, twist-3 and twist-4 contributions corresponding to the first, second, and third terms inside the brackets in the right hand side of Eq. \#11. Only the sum rules for \( \Lambda_{10}^{D^*\bar{D}^*} \) up to \( \Lambda_{14}^{D^*\bar{D}^*} \) get only twist-4 contributions, and give results compatible with zero. It is also interesting to notice that if we consider only the twist-2 contributions we recover the sum rules obtained in ref. \#12.

The results for all other sum rules show a similar behavior and the amplitude can be extracted by the extrapolation of the fit to \( M^2 = 0 \). The values for all the parameters are given in \#23. In Eq. \#12 the structures multiplying \( \Lambda_4 \) and \( \Lambda_5 \) break chiral symmetry \#14 and, therefore, will be neglected.

Having the QCD sum rule results for the amplitude of the three processes \( J/\psi \pi \rightarrow \bar{D} \bar{D}^*, \bar{D} D, \bar{D}^* D^* \), we can evaluate the cross section. In Fig. 2 we show the cross section for the \( J/\psi - \pi \) dissociation. The shaded area give an evaluation of the uncertainties in our calculation obtained with the two procedures described above. It is important to keep in mind that, since our sum rule was derived in the limit \( p_1 < p_2, p_3, p_4 \), we can not extend our results to large values of \( \sqrt{s} \).
FIG. 2: Total $J/\psi \pi$ dissociation cross sections of the processes $J/\psi \pi \rightarrow \bar{D} D^* + D D^* + \bar{D} D + D^* D^*$. The shaded area give an evaluation of the uncertainties in our calculation.

In a hadron gas, pions collide with the $J/\psi$ at different energies. The momentum distribution of thermal pions in a hadron gas depends on the effective temperature $T$ with an approximate Bose-Einstein distribution. Therefore, the relevant quantity is not the value of the cross section at a given energy, but $\langle \sigma \pi J/\psi \rangle$ which is the product of the dissociation cross section and the relative velocity averaged over the energies of the pions.

As shown in Fig. 3, $\langle \sigma \pi J/\psi \rangle$ increases with the temperature. Since the $J/\psi$ dissociation by a pion requires energetic pions to overcome the energy threshold, it has a small thermal average at low temperatures. The shaded area in Fig. 3 give an evaluation of the uncertainties in our calculation due to the two procedures used to extract the hadronic amplitudes.

In conclusion, we have studied the $J/\psi$ dissociation cross section by pions using the QCDSR technique, based on a three-point function using vacuum-pion correlation functions. We have estimated the hadronic amplitudes by working up to twist-4 in the limit $p_1 << p_2, p_3, p_4$. Our results are in agreement with the former QCDSR calculation, done with a four-point function at the pion pole \cite{15}. Our results for the cross section as a function of $\sqrt{s}$ are smaller than the results using meson-exchange models (without form factors), but larger than the calculation based on quark-exchange models.

The dominant contribution to the hadronic amplitudes comes from the twist-2 operator, or equivalently, from the quark condensate. As we know that the quark condensate is stronger in the vacuum and weaker in the interior of hadrons, we can conclude that the charmronium “sees” and interacts with the surface of the pions, where there is a “halo” of condensates. This is way the cross section can be larger than the geometric value. In our approach the continuous growth of the cross section comes from the growth of the phase space, as in the effective Lagrangian calculations. In the short distance QCD calculation \cite{4,10} the cross section also grows with $\sqrt{s}$, but the growth there is considerably smaller because the rise in the gluon density cannot completely compensate the fall of the partonic cross section \cite{10}.

The thermal average of the $J/\psi - \pi$ dissociation cross section increases with the temperature and at $T = 150$ MeV we get $\langle \sigma \pi J/\psi \rangle \sim 0.2 - 0.4$ mb which is smaller than the values used in phenomenological studies of $J/\psi$ absorption by comoving hadrons in relativistic heavy ion collisions.

The same approach used here could be applied to calculate the $\Upsilon \pi$ cross section. In a recent work \cite{10} the $\sigma_{\Upsilon \pi}$ was computed using short distance QCD. Since we expect the non-perturbative corrections to be less important for heavier systems, the differences between short distance QCD and QCD sum rules should be smaller for the $\Upsilon \pi$ system, and a systematic comparison between the two approaches becomes possible. We will address this point in the future.

Another possible extension of this work is the calculation of the $\chi_c \pi$ and $\eta_c \pi$ cross sections. This can be done by replacing the $j_{\mu}^0$ current in Eq. (1) by the corresponding $\chi_c$ and $\eta_c$ currents. Unfortunately since $\psi'$ and $J/\psi$ have the same quantum numbers and are, therefore, described by the same current, it is not possible, in this approach, to estimate the $\psi' \pi$ cross section. The $\psi'$ contribution to the present sum rule calculation is inside the parameters $A$ and $B$, in Eq. (11) and cannot be separated from the other higher mass states contributions.
We are grateful to J. Hufner and H. Kim for fruitful discussions. M.N. would like to thank the hospitality and financial support from the Yonsei University during her stay in Korea. This work was supported by CNPq and FAPESP-Brazil.

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