Effect of in-medium spectral density of $D$ and $D^*$ mesons on the $J/\psi$ dissociation in hadronic matter

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Abstract

The one-loop self-energy of the $D$ and $D^*$ mesons in a hot hadronic medium is evaluated using the real time formalism of thermal field theory. The interaction of the heavy open-charm mesons with the thermalized constituents ($\pi, K, \eta$) of the hadronic matter is treated in the covariant formalism of heavy meson chiral perturbation theory. The imaginary parts are extracted from the discontinuities of the self-energy function across the unitary and the Landau cuts. The non-zero contribution from the latter to the spectral density of $D$ and $D^*$ mesons opens a number of subthreshold decay channels of the $J/\psi$ leading to a significant increase in the dissociation width in hadronic matter.

1 Introduction

Heavy quarks, charm and bottom, and their bound states have emerged in a leading role as probes of the strongly interacting system produced in heavy ion collisions [1, 2]. In fact, along with the elliptic flow and nuclear suppression of light hadrons, these quantities have also been measured for the single $e^-$ spectra from open charm and beauty meson decays at RHIC [3, 4]. After the discovery of QGP as a near perfect fluid at RHIC, much attention has gone into the study of its transport properties and this is where heavy quarks are advantageous. Because of the fact that their masses are much larger than the typical temperature even at LHC energies, they are predominantly produced in the early primordial stages of the collision and are witness to the entire evolution. Moreover, their thermalisation in the QGP being unlikely they presumably retain the memory of their interaction. Their large mass allows them to be treated as Brownian particles validating a Fokker-Planck description for the study of transport properties using heavy quarks as probes. This formalism has been implemented in many analyses [5, 7, 6, 8] and the drag and diffusion coefficients were evaluated.

In most estimates the role of hadronic matter in the analysis of elliptic flow and nuclear suppression of heavy flavors was completely ignored. Recently, the drag and diffusion coefficients of a hot hadronic medium consisting of pions, kaons and eta using $D$ [9] and $B$ [10] mesons as well as their role in heavy flavour suppression [11] were evaluated using interactions from heavy meson chiral perturbation theory. The significant values obtained for these coefficients indicate a substantial interaction of the heavy mesons with the thermal bath (see also [12, 13, 14]). It is hence worthwhile to investigate how the same interaction modifies the spectral properties of $D$ and $D^*$ mesons in hot hadronic matter. This in turn should be reflected in observables like dissociation width of the $J/\psi$.
and the dilepton spectra in the intermediate mass region where correlated $D$ meson decays make a significant contribution $^{15}$. The suppression of the $J/\psi$ production in highly relativistic heavy ion collisions with respect to binary scaled $pp$ collisions which is experimentally found by NA50 $^{16}$, NA60 $^{17}$ as well as by the PHENIX $^{18}$ collaboration has long been suggested as a signal of quark-gluon plasma formation in heavy ion collisions $^{19, 20}$. However, other mechanisms based on the $J/\psi$ absorption by comoving hadrons have also been proposed as a possible explanation $^{1}$. In this connection the inelastic reaction rates of $J/\psi$ in the hadronic phase have been studied using different effective hadronic models $^{21, 22, 23}$. In addition, several theoretical efforts $^{24, 25, 26, 27, 28, 29, 30, 31, 32, 33, 34, 35, 36, 37}$ have been made in order to understand the in-medium behavior of charmed mesons using different hadronic models e.g. QCD sum rules $^{25, 26, 27}$, Quark-Meson Coupling (QMC) $^{28}$, Nambu-Jona-Lasinio (NJL) model $^{29, 30}$. Most of them have predicted a larger mass drop of open-charm mesons than that of $J/\psi$ which may help to explain the $J/\psi$ suppression in a hadronic environment $^{24, 25, 32}$ although a width enhancement with negligible mass shift of open charm mesons $^{31, 30, 34}$ have also been suggested in order to explain the suppression $^{30, 34}$.

In the present work we investigate in detail the in-medium modification of open charm mesons $D$ and $D^*$ by performing an explicit evaluation of their spectral functions in the Real-time formalism of Thermal Field theory (RTF). The interaction vertices appearing in the one-loop self-energy of $D$ and $D^*$ is obtained from the Lagrangian of Chiral Perturbation Theory ($\chi PT$) involving heavy-light pseudoscalar and vector mesons in the covariant form $^{38}$. Unlike most treatments in the literature, we follow the technique of extracting the imaginary part from the discontinuities of the self-energy function in the complex energy plane $^{39, 40}$. The position of the branch cuts, the Landau cut in particular, which appears only in the medium along with the usual unitary cut which also exists in vacuum but is now weighted by the phase space distributions, are determined from the kinematics before evaluating their relative contributions. Using the in-medium spectral function of $D$ and $D^*$ mesons we then investigate its effect on the $J/\psi$ dissociation in hadronic matter by calculating its width in the different $D\overline{D}$, $D^*\overline{D}$ and $D^*\overline{D}$ channels.

The manuscript is organized as follows. Section II deals with the self-energy of $D$ and $D^*$ mesons and begins with a discussion of the structure of the propagator in the RTF. After a discussion of the interaction vertices obtained from the chiral Lagrangian, the branch cut structure of the one-loop self-energy is elaborated upon. The results are presented in section III, followed by a summary in section IV.

2 The self-energy of the $D$ and $D^*$

2.1 Structure of interacting propagator in RTF

In the real time formalism of thermal field theory $^{41}$, all two point functions have a $2\times2$ matrix structure. These matrices may however be diagonalized in terms of a single analytic function that determines completely the dynamics of the corresponding two point function $^{42, 43}$. Since this function is related simply to any one component, say the $11$ component of the matrix, we need to calculate only this component of the matrix.

The $11$ component of the free thermal matrix propagator has two parts. Along with the usual vacuum term there is a part containing the on-shell distribution functions of like particles in the medium. Here pseudoscalar and vector mesons will appear in our calculation and hence we discuss their $11$ components of the propagators in the following. For the pseudoscalar, this is given by

$$G^{(0)11}(q) = \Delta^{(0)}(q) + 2\pi i\delta(q^2 - m_\pi^2)$$

(1)
where $\Delta^{(0)}$ is the vacuum propagator of the scalar field and $n$ its equilibrium distribution function given by

$$
\Delta^{(0)}(q) = \frac{-1}{q^2 - m^2 + i\eta}, \quad n(\omega_q) = \frac{1}{e^{\beta \omega_q} - 1}
$$

with $\omega_q = \sqrt{q^2 + m^2}$.

The Dyson equation relates the complete propagator (matrix) $G^{ab}$ with the free propagator $G^{(0)ab}$ and self-energy $\Pi^{ab}$ matrices at finite temperature,

$$
G^{ab} = G^{(0)ab} - G^{(0)ac} \Pi^{cd} G^{db}.
$$

where $a, b, c, d$ are thermal indices and take the values 1 and 2. These matrices can be diagonalized in terms of their respective analytic functions (denoted by a bar) which again obey

$$
\overline{G} = \overline{G}^{(0)} - \overline{G}^{(0)} \overline{\Pi} \overline{G}, \text{ where } \overline{G}^{(0)} = \Delta^{(0)}(q)
$$

and can be solved to get the full propagator $\overline{G}$. The diagonal element of self-energy function is related to the 11 component as

$$
\text{Re}\Pi(q) = \text{Re}\Pi^{11}(q)
$$

$$
\text{Im}\Pi(q) = \epsilon(q_0) \tanh(\frac{\beta q_0^2}{2}) \text{Im}\Pi^{11}(q)
$$

in terms of which the spectral function of the $D$ meson is defined as

$$
A_D = 2\epsilon(q_0) \text{Im}\overline{G}(q)
$$

where

$$
\text{Im}\overline{G}(q) = \frac{-\sum \text{Im}\Pi}{(q^2 - m^2_D - \sum \text{Re}\Pi)^2 + (\sum \text{Im}\Pi)^2},
$$

the summation running over all the loops. To obtain the in-medium spectral function one thus evaluates the real and imaginary parts of the 11-component of the self-energy.

The $D^*$ being a vector meson, the propagator contains tensor indices. The 11-component in this case is given by,

$$
G^{(0)11}_{\mu\nu}(p) = (-g_{\mu\nu} + \frac{p_{\mu} p_{\nu}}{m_p^2}) G^{(0)11}(p)
$$

As shown in [42, 43] the full propagator can be obtained analogously as above. The spectral function is then given by eq. (7) above with the spin-averaged self-energy

$$
\overline{\Pi} = \frac{1}{3} \overline{\Pi}^\mu.
$$

### 2.2 Interaction vertices

The vertices appearing in the one-loop self-energy graphs are obtained using chiral perturbation theory. The lowest order chiral Lagrangian for the heavy-light pseudoscalar and vector meson is [38],

$$
\mathcal{L} = \langle D_\mu PD^{\mu} P^\dagger \rangle - m_D^2 \langle PP^\dagger \rangle - \langle D_\mu P^{*\nu} D^{\mu} P^{*\dagger}_{\nu} \rangle + m_D^2 \langle P^{*\nu} P^{*\dagger}_{\nu} \rangle + \frac{g}{2m_D} \langle (P^{*\mu} P^{*\dagger}_{\mu} - P^{*\mu} P^{*\dagger}_{\mu}\gamma_5) e^{i\omega_0^\beta} \rangle
$$
Figure 1: One loop diagram of $D$ meson self-energy where $\Phi$ stands for $\pi, \eta$ and $K$ mesons

Figure 2: One loop diagram of $D^*$ meson self-energy

where $P = (D^0, D^+, D^+_s)$ and $P^*_\mu = (D^{*0}, D^{*+}, D^{*_s})_\mu$ are the triplets of $D$ and $D^*$ meson fields. The covariant derivatives are defined as $D_\mu P_a = \partial_\mu P_a - \Gamma^b_{a\mu} P_b$ and $D^\mu P^*_a = \partial^\mu P^*_a + \Gamma^b_{ab} P^*_b$, with $a, b$ the $SU(3)$ flavor indices. $m_D$ is the mass of heavy-light meson in the chiral limit. The value of the heavy-light pseudoscalar-vector coupling constant $g = 1.177$ GeV is obtained by reproducing the experimental $D^* \to D\pi$ decay width of $\sim 65$ keV with the above interaction\[38\]. The vector and axial-vector currents are respectively given by

$$\Gamma_\mu = \frac{1}{2} (u^\dagger \partial_\mu u + u \partial_\mu u^\dagger)$$
$$u_\mu = i (u^\dagger \partial_\mu u - u \partial_\mu u^\dagger) \tag{11}$$

where $u = \exp(i \lambda_\alpha \Phi_\alpha)$ collects the octet of Nambu Goldstone fields with

$$\lambda_\alpha \Phi_\alpha = \sqrt{2} \left( \begin{array}{cccc} \frac{\pi^0}{\sqrt{2}} & \eta & K^+ & K^- \\ \frac{\pi^+}{\sqrt{2}} & \frac{\pi^-}{\sqrt{6}} & -\frac{\eta}{\sqrt{6}} & -2\eta \\ K^0 & \bar{K}^0 & \frac{\eta}{\sqrt{6}} & \frac{\eta}{\sqrt{2}} \\ \frac{2}{\sqrt{6}} \eta & \frac{2}{\sqrt{2}} \eta & \frac{3}{2} \eta & \frac{3}{2} \eta \\ \end{array} \right) \tag{12}$$

and $f_\pi$ is the pseudoscalar decay constant in the chiral limit. To lowest order in $\Phi$ the vector and axial-vector currents are

$$\Gamma_\mu = \frac{1}{8 f^2_\pi} [\Phi, \partial_\mu \Phi], \quad u_\mu = -\frac{1}{f_\pi} \partial_\mu \Phi \tag{13}$$

from which the interaction terms are separately obtained as

$$L_{P \Phi P \Phi} = \frac{1}{8 f^2_\pi} (\partial_\mu P^\dagger \Phi, \partial^\mu \Phi) P^\dagger - P[\Phi, \partial^\mu \Phi] \partial_\mu P^\dagger \tag{14}$$

$$L_{P^* \Phi P^{*\dagger} \Phi} = -\frac{1}{8 f^2_\pi} (\partial_\mu P^{*\dagger} \Phi, \partial^\mu \Phi) P^{*\dagger} - P^{*\dagger}[\Phi, \partial^\mu \Phi] \partial_\mu P^{*\dagger} \tag{15}$$


\[ L_{\nu\gamma P_{\mu}^{\nu}} = -\frac{g}{f_{\pi}} (P_{\mu}^{\nu} \partial^\mu \Phi P_{\mu}^{\nu} - P \partial^\mu \Phi P_{\mu}^{\nu}) \]  
\[ L_{\nu\gamma P_{\mu}^{\nu}} = -\frac{g}{2m_{D}\pi} \langle (P_{\mu}^{\nu} \partial_\alpha \Phi \partial_\beta P_{\mu}^{\nu} - \partial_\beta P_{\mu}^{\nu} \partial_\alpha \Phi P_{\mu}^{\nu}) e^{\mu\nu\alpha\beta} \rangle \]  

Considering e.g. the \( D^+ \) from the triplet \( P \), the relevant part of Eq.\,(14) and \,(16) are respectively given by

\[ L_{D^+ P_{\mu}^{\nu} D^+} = -\frac{1}{4f_{\pi}^2} ((D^+ \partial_\mu \pi^0 \partial^\mu D^+ - \partial^\mu D^+ \partial_\mu \pi^0 + \partial^\mu D^+ \pi D^+ - \partial^\mu \pi D^+ \partial_\mu \pi^0) \]  
\[ + (D^+ \partial_\mu \vec{K}_0 \partial^\mu \vec{K} - \partial^\mu D^+ \partial_\mu \vec{K}_0 \vec{K} + \partial^\mu D^+ \vec{K}_0 \vec{K} D^+ - \vec{K} \partial^\mu \partial_\mu \vec{K}_0 \vec{K}) \]  

\[ L_{D^+ P_{\mu}^{\nu}D^+} = \frac{i g}{f_{\pi}} \sqrt{2(\partial^\mu \pi^0 D^+ - \partial^\mu \pi D^+) + \sqrt{2(\partial^\mu \pi D^+ - \partial^\mu \pi D^+ K^0 - \partial^\mu \pi D^+ \vec{K} D^+ - \partial^\mu \pi D^+ \vec{K} K^0 \vec{K})} \]  

Similarly for \( D^{++} \) from the triplet \( P^* \) the relevant part of Eq.\,(15), \,(16) and \,(17) are respectively given by

\[ L_{D^{++} P_{\mu}^{\nu} D^{++}} = \frac{1}{4f_{\pi}^2} \left[(D^{++} \partial_\mu \pi^0 \partial^\mu D^{++} - \partial^\mu D^{++} \partial_\mu \pi^0 + \partial^\mu D^{++} \pi D^{++} - \partial^\mu \pi D^{++} \partial_\mu \pi^0) \right. 
\[ - (D^{++} \partial_\mu \vec{K}_0 \partial^\mu \vec{K} - \partial^\mu D^{++} \partial_\mu \vec{K}_0 \vec{K} + \partial^\mu D^{++} \vec{K}_0 \vec{K} D^{++} - \vec{K} \partial^\mu \partial_\mu \vec{K}_0 \vec{K}) \]  

\[ L_{D^{++} P_{\mu}^{\nu}D^{++}} = \frac{i g}{f_{\pi}} \sqrt{2(\partial^\mu \pi^0 D^{++} - \partial^\mu \pi D^{++}) + \sqrt{2(\partial^\mu \pi D^{++} - \partial^\mu \pi D^{++} K^0 - \partial^\mu \pi D^{++} \vec{K} D^{++} - \partial^\mu \pi D^{++} \vec{K} K^0 \vec{K})} \]  

\[ L_{D^{++} P_{\mu}^{\nu}D^{++}} = \frac{g e^{\mu\nu\alpha\beta}}{2m_{D}\pi} \sqrt{2(\partial_\beta D^{++} \partial_\alpha \pi^0 \partial^\mu \vec{K} - D^{++} \partial_\alpha \pi^0 \partial_\beta \vec{K} D^{++} - \partial_\beta D^{++} \partial_\alpha \pi^0 \partial_\mu \vec{K} D^{++} \vec{K} + \partial^\mu \partial_\alpha \pi^0 \partial_\beta \vec{K} \vec{K} \partial_\mu \partial_\alpha \pi^0 \partial_\beta \vec{K} \vec{K})} \]  

Having fixed the propagator and interaction vertices we are now in a position to write down the expression for the self-energy function. The expression of \( D^+ \) meson self-energy coming from fig\,(1A) and (B) are respectively given by

\[ \Pi^{11}(q) = i \int \frac{d^4k}{(2\pi)^4} N(q,k)D^{11}(k,m_k)D^{11}(q-k,m_{\pi}) \]  
\[ \Pi^{11}(q) = i \int \frac{d^4k}{(2\pi)^4} N(q,k)D^{11}(k,m_k) . \]  

The \( D^{++} \) being a vector meson its self-energy will have tensor indices. The corresponding expressions for fig.\,(2A), \,(2B) and \,(2C) are given by Eqs.\,(23) and \,(24) with \( N(q,k) \to N^{\mu\nu}(q,k) \) and consequently \( \Pi^{11}(q) \to \Pi^{\mu\nu,11}(q) \). The term \( N(q,k) \) \((and \, N^{\mu\nu}(q,k)) \) results from factors coming from the vertices and propagators in the self-energy diagrams.
We have considered four possible loops $\pi^+D^0$, $\pi^0D^{*+}$, $\eta D^*$ and $K^0D^{*+}$ to evaluate the $D^+$ and $D^{*+}$ meson self-energy in the medium which are represented in fig.(1A) and fig.(2B) respectively. They appear due to the interactions in (19) and (22) respectively. The $\pi^+D^0$, $\pi^0D^+$, $\eta D^*$ and $K^0D^{*+}$ loops additionally appear for the $D^{*+}$ meson self-energy due to the interactions in (21) which are shown in fig.(2A). The form of $N(q,k)$ and $N^{\mu\nu}(q,k)$ in the expression of $D$ and $D^*$ meson self-energy are respectively obtained as

$$N(k,q) = -\alpha^2\left(\frac{g}{f_{\pi}}\right)^2[k^2 - \frac{(k \cdot q - q^2)^2}{m_p^2}]$$ [for fig.(1A), from Eq.(19)]

$$N^{\mu\nu}(k,q) = -\alpha^2\left(\frac{g}{m_D f_{\pi}}\right)^2[k^2 q^2 A^{\mu\nu} + B^{\mu\nu}]$$ [for fig.(2B), from Eq.(22)]

where $A^{\mu\nu} = -g^{\mu\nu} + \frac{q^\mu q^\nu}{q^2}$, $B^{\mu\nu} = q^2 k^\mu k^\nu - (q \cdot k)(q^\mu k^\nu + q^\nu k^\mu) + (q \cdot k)^2 g^{\mu\nu}$.

In above equations, $\alpha = \sqrt{2}, 1, \sqrt{2}, \frac{1}{\sqrt{3}}$ for $\pi^+D^0$ ($\pi^+D^0$), $\pi^0D^{*+}$ ($\pi^0D^{*+}$), $K^0D^{*+}$ ($K^0D^{*+}$) and $\eta D^{*+}$ ($\eta D^{*+}$) loops respectively and $m_p$ in Eq.(21) is the mass of $D^*$ meson. Using the interaction Lagrangian (18) for $D^+$ and (20) for $D^{*+}$, the integrands in Eq.(24) turn out to be odd functions of the integration variable $k$. Therefore fig.(1B) and fig.(2C) do not contribute to the spectral function.

### 2.3 Branch cut structure of the self energy

The integral over $k_0$ is easily performed by choosing suitable contours. The imaginary part of the self-energy function is then obtained from the 11-component using (4) to get

$$\text{Im} \Pi(q) = -\pi \int \frac{d^3 \bar{k}}{(2\pi)^3 4\omega_{k}\omega_{p}} \times$$

$$[N(k_0 = \omega_k)\{(1 + n(\omega_k) + n(\omega_p))\delta(q_0 - \omega_k - \omega_p)$$

$$+ (n(\omega_k) - n(\omega_p))\delta(q_0 - \omega_k + \omega_p)\}$$

$$+ N(k_0 = -\omega_k)\{(n(\omega_k) - n(\omega_p))\delta(q_0 + \omega_k - \omega_p)$$

$$- (1 + n(\omega_k) + n(\omega_p))\delta(q_0 + \omega_k + \omega_p)\}]$$

(28)

where $\omega_k = \sqrt{k^2 + m_k^2}$ and $\omega_p = \sqrt{(q - \bar{k})^2 + m_p^2}$ are the energies of light pseudoscalar mesons and the heavy-light charm mesons respectively. The cuts in the self-energy function correspond to regions in which the four terms in the above expression are non-vanishing and are determined by the respective $\delta$-functions. A little inspection suggests that the first and the fourth terms are finite for $q^2 \geq (m_p + m_h)^2$ giving rise to the unitary cut whereas the second and third terms are non-vanishing for $q^2 \leq (m_p - m_h)^2$ producing the Landau cut. The unitary cut arises as a consequence of gain or loss of $D$ (or $D^*$) mesons due to formation from or decay into the particles in the medium. The Landau type of discontinuity appears due to disappearance of the heavy mesons as a result of scattering with the bath particles. While the unitary cut exists in the vacuum also, the Landau cut is purely a medium effect.
The thermal contribution to the real part is given by
\[
\text{Re} \Pi(q) = \mathcal{P} \int \frac{d^3q}{(2\pi)^3} n(\omega_k) \left\{ \frac{N(k_0 = \omega_k)}{(q_0 - \omega_k)^2 - \omega_p^2} + \frac{N(k_0 = -\omega_k)}{(q_0 + \omega_k)^2 - \omega_k^2} + \frac{n(\omega_p)}{2\omega_p} \left( \frac{N(k_0 = q_0 + \omega_p)}{(q_0 - \omega_p)^2 - \omega_k^2} + \frac{N(k_0 = q_0 - \omega_p)}{(q_0 + \omega_p)^2 - \omega_k^2} \right) \right\}
\]
(29)
where \(\mathcal{P}\) indicates the principal value. The vacuum contribution is divergent and is assumed to renormalize the mass of the \(D\) mesons to their physical values.

Writing \(d^3q = 2\pi \sqrt{\omega_k^2 - m_k^2} \omega_k d\omega_k \sin \theta d\theta d\phi\), where \(\theta\) is the angle between \(\vec{q}\) and \(\vec{k}\), we can readily integrate over \(\cos \theta\) using the \(\delta\)-functions yielding
\[
\cos \theta_0 = \frac{-R^2 \pm 2q_0\omega_k}{2|q|\sqrt{\omega_k^2 - m_k^2}}, \quad R^2 = q^2 - m_p^2 + m_\pi^2.
\]
(30)
with ‘+’ the for first two terms and ‘−’ for the last two terms of Eq. (28).

In the present study we will be interested in the region \(q^2 > 0\) and \(q_0 > 0\). The imaginary parts are given by [40]
\[
\text{Im} \Pi = -\frac{1}{16\pi|q|} \int_{\omega_-}^{\omega_+} d\omega N(k_0 = -\omega_k)
\]
(31)
\[
\{ n(\omega_k) - n(\omega_k + q_0) \}
\]
\[
\text{Im} \Pi = -\frac{1}{16\pi|q|} \int_{\omega_-}^{\omega_+} d\omega N(k_0 = \omega_k)
\]
(32)
\[
\{ 1 + n(\omega_k + q_0) + n(\omega_k) \}
\]
for \(|q| \leq q_0 \leq \sqrt{(m_p - m_k)^2 + |q|^2}\) (region of Landau cut) and
\[
\text{Im} \Pi = -\frac{1}{16\pi|q|} \int_{\omega_-}^{\omega_+} d\omega N(k_0 = \omega_k)
\]
\[
\{ 1 + n(\omega_k + q_0) + n(\omega_k) \}
\]
for \(q_0 \geq \sqrt{(m_p + m_k)^2 + |q|^2}\) (region of unitary cut) with \(\omega_\pm = \frac{R^2}{2q^2}(q_0 \pm |q|v)\), \(\omega_{\pm} = \frac{R^2}{2q^2}(-q_0 \mp |q|v)\) and \(v(q^2) = \sqrt{1 - \frac{4m_p^2m_\pi^2}{R^4}}\).

3 Results and Discussion

We now present results of numerical calculation. For the \(D\) meson \(m_k = m_\pi, m_\eta, m_K\) and \(m_p = m_D, m_{D^*}\) whereas for the \(D^*\) meson \(m_p = m_D, m_{D^*}\). The \(\text{Im} \Pi\) (upper panels) and \(\text{Re} \Pi\) (lower panels) of \(D\) meson for \(D^*\Phi\) loops is shown in fig. (3) whereas same of \(D^*\) meson are shown in fig. (4) and (5) for \(D\Phi\) and \(D^*\Phi\) loops respectively. In the upper panels of both curves, the branch cut structure of \(\Pi(q^0, q)\) is clearly seen. The Landau cut contributes in the region \(0 \leq M \leq (m_p - m_k)\) and the unitary cut for \(M \geq (m_p + m_k)\) thus leaving in-between a region where the imaginary part is exactly zero. It is easy to see that for the \(D\) meson which has a nominal mass of 1867 MeV, only the Landau cut of the \(D^*\pi\) loop contributes at the pole whereas for the \(D^*\) meson the the unitary cut of the \(D\pi\) loop contributes at the \(D^*\) pole (2008 MeV). The real part however receives contribution from all the four terms in Eq. (29) as seen in the lower panels of all these figures.

We now present the results for the spectral function of the \(D\) and \(D^*\) mesons. For the vector case we take a spin average of the imaginary part of the full propagator as done for the self-energy in
Figure 3: Imaginary and real part of $D$ meson self-energy for different $D^*\Phi$ loops in the upper and lower panel respectively.

Figure 4: Imaginary and real part of $D^*$ meson self-energy for different $D\Phi$ loops in the upper and lower panel respectively.

Figure 5: Imaginary and real part of $D^*$ meson self-energy for different $D^*\Phi$ loops in the upper and lower panel respectively.
Figure 6: Landau part of spectral function of $D$ and $D^*$ in the upper and lower panel respectively for three different temperatures.

Figure 7: The unitary part of the spectral functions of $D$ and $D^*$ mesons in the upper and lower panels respectively at $T = 175$ MeV.
the expression for the in-medium \( J/\psi \) where
\[
\vec{P}_{cm} = \sqrt{\left| m_\psi - (p_D + p_{\bar{D}})^2 \right|} \frac{m_\psi}{2m_\psi}
\]
is the center of mass momentum and \( p_D = \sqrt{(p_D^0)^2 - p_{D\perp}^2} \)
is the invariant masses of \( D \) and \( \bar{D} \) respectively. We take \( g_\psi = 7.8 \) as in [24].

Figure 8: The total dissociation width of \( J/\psi \) plotted against temperature.

\[\text{eq. (34)}\]
Since the entire spectral shape is of importance and not only its value at the pole, we show the contributions from the Landau and unitary parts separately in figs. [6] and [7] respectively. As discussed previously the strength in the Landau region stems from processes in which the dissipative interactions with the thermalised light mesons \( \pi, K \) and \( \eta \). Analogously decay and regeneration processes involving the \( D \) and \( D^* \) mesons leads to the strength in the unitary region as shown in fig. [7].

The observed multiple structures in the spectral function have a non-trivial effect on the \( J/\psi \) dissociation in hadronic matter. For this purpose we treat the \( D \) and \( D^* \) mesons mesons as resonances with temperature dependent finite collisional widths as described by their spectral densities. With the help of the phenomenological Lagrangian densities [46] [22]

\[
\mathcal{L}_{\psi DD} = ig_\psi \psi^\mu [\vec{D} (\partial_\mu D) - (\partial_\mu \vec{D}) D],
\]
\[
\mathcal{L}_{\psi DD^*} = \frac{g_\psi}{m_\psi} \epsilon_{\alpha \beta \mu \nu} (\partial^\alpha \psi^\beta \Delta \mu \nu D + \Delta \mu \nu D) ,
\]
\[
\mathcal{L}_{\psi D^* D^*} = ig_\psi \{ (\partial_\mu \vec{D}^{\mu}) D^*_{\nu} - \vec{D}^{\nu} (\partial_\mu D^*_{\nu}) \} + \{ (\partial_\mu \psi^\nu) \Delta_{\mu \nu} \}
\]

the expression for the in-medium \( J/\psi \) dissociation width for \( DD \), \( DD^* \) and \( D^* D^* \) channels are given by,

\[
\Gamma_{\text{med}}(J/\psi \rightarrow DD) = \int \frac{g_\psi^2 F^2}{3m_\psi^2} \left| \vec{P}_{cm} (p_D^2, p_{\bar{D}}^2) \right|^3 \frac{Im G_D(p_D^0, \vec{P}_{cm})}{\pi} \frac{Im G_{\bar{D}}(p_{\bar{D}}^0, \vec{P}_{cm})}{\pi} \, dp_D^2 \, dp_{\bar{D}}^2
\]

\[
\Gamma_{\text{med}}(J/\psi \rightarrow DD^*) = \int \frac{g_\psi^2 F^2}{3m_\psi^2} \left| \vec{P}_{cm} (p_D^2, p_{D^*}^2) \right|^3 \frac{Im G_D(p_D^0, \vec{P}_{cm})}{\pi} \frac{Im G_{D^*}(p_{D^*}^0, \vec{P}_{cm})}{\pi} \, dp_D^2 \, dp_{D^*}^2
\]

\[
\Gamma_{\text{med}}(J/\psi \rightarrow D^* D^*) = \int \frac{g_\psi^2 F^2}{3m_\psi^2} \left| \vec{P}_{cm} (p_{D^*}^2, p_{D^*}^2) \right|^3 \frac{Im G_D(p_{D^*}^0, \vec{P}_{cm})}{\pi} \frac{Im G_{D^*}(p_{D^*}^0, \vec{P}_{cm})}{\pi} \, dp_{D^*}^2 \, dp_{D^*}^2
\]
After integrating over $p_{D}^{2}$ and $p_{D^*}^{2}$ the total $J/ψ$ dissociation width for all the channels ($DD$, $DD^*$ as well as $D^*D^*$) is shown in fig. (8) as a function of temperature. In order to effectively account for the composite nature of the hadrons involved we incorporate a form factor, $F = \frac{\Lambda^2}{\Lambda^2 + p^2_{cm}}$ at the vertex and show results for two values of the cut-off $\Lambda$. The solid line for which $\Lambda=1$ GeV appears to be in reasonable agreement with [45].

The above discussions indicate that the opening up of subthreshold decay channels of $J/Ψ$ as shown here is primarily on account of the broadening of the spectral shape of the charmed mesons and not so much due to the shift of their pole positions towards lower masses. It is nevertheless worthwhile to investigate the role of the effective mass of the $J/Ψ$. We thus evaluate its self-energy following the same formalism (as for the $D^*$ meson) for $D\bar{D}$, $DD^*$ and $D^*D^*$ loops. The spin averaged real part is given by eq. (29) where the expressions for $N(q,k)$ for the three cases are given by

$$N^{DD}(q,k) = \frac{2g_{\psi}^2 F^2}{3}(q^2 + 4k^2 - 4k \cdot q)$$

$$N^{DD^*}(q,k) = \frac{8g_{\psi}^2 F^2}{3m_{\psi}^2}(q^2 k^2 - (k \cdot q)^2)$$

$$N^{D^*D^*}(q,k) = \frac{2g_{\psi}^2 F^2}{3} \left[ 18(k^2 - k \cdot q + q^2) + \frac{q^4}{m_D^4}(k^2 q^2 - (k \cdot q)^2) - \frac{1}{m_D^4}(6k^4 - 12k^2 q \cdot q + 8(k \cdot q)^2) - 2q^2 k^2 + 3q^4 \right].$$

In fig. (9) we plot the mass shift of the $J/Ψ$, $\Delta_{\psi} = \sqrt{m_{\psi}^2 + \text{Re}\Pi} - m_{\psi}$ as a function of its invariant mass at $T=200$ MeV. The shift turns out to be negligible at the pole.

4 Summary and Conclusion

In summary, motivated by the non-trivial drag and diffusion coefficients of $D$ mesons in hadronic matter seen in [9], [10] and [11] we have performed a detailed analysis of the spectral properties of heavy open charm mesons in hadronic matter at finite temperature. We use the real-time formalism of thermal field theory and extract the imaginary part of the self-energy from the discontinuities.
in the complex energy plane. We obtain a small spectral broadening of the charm mesons at finite
temperature which enables the $J/\psi$ to decay into the subthreshold $D\bar{D}$, $D^*\bar{D}$ as well as $D^*\bar{D}$
channels thus contributing to its suppression in hadronic environment. We also find that the mass
of the $J/\Psi$ does not change appreciably in this model. The magnitude of the real part and the
decay width however depend crucially on the couplings as well as the cut-off employed.

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References

[1] R. Vogt, Phys. Rep. 310, 197 (1999).

[2] R. Rapp and H. van Hees, R. C. Hwa, X.-N. Wang (Ed.) Quark Gluon Plasma 4, World
Scientific, 111 (2010)

[3] S. S. Adler et al. (PHENIX Collaboration), Phys. Rev. Lett. 96, 032301 (2006); B. I. Abeleb
et al. (STAR Collaboration), Phys. Rev. Lett. 98, 192301 (2007).

[4] S. S. Adler et al. (PHENIX Collaboration), Phys. Rev. Lett. 98, 172301 (2007).

[5] B. Svetitsky, Phys. Rev. D 37 (1988) 2484.

[6] M. G. Mustafa, D. Pal and D. Kumar Srivastava, Phys. Rev. C 57 (1998) 889 [Erratum-ibid.
C 57 (1998) 3499]

[7] H. van Hees and R. Rapp, Phys. Rev. C 71 (2005) 034907

[8] S.K. Das, J. Alam, and P. Mohanty, Phys. Rev. C 80, 054916 (2009); S.K. Das, J. Alam, P.
Mohanty, and B. Sinha, Phys. Rev. C 81, 044912 (2010)

[9] S. Ghosh, S.K. Das , S. Sarkar and J. Alam, Phys. Rev. D 84, 011503 (2011)

[10] S.K. Das, S. Ghosh, S. Sarkar, J. Alam Phys. Rev. D 85 (2012) 074017

[11] S.K. Das, S. Ghosh, S. Sarkar, J. Alam Phys. Rev. D 88 (2013) 017501

[12] M. Laine, JHEP 04, 124 (2011).

[13] M. He, R. J. Fries and R. Rapp, Phys. Let. B 701, 445 (2012)

[14] L. Abreu, D. Cabrera, F. J. Llanes-Estrada and J. M. Torres-Rincon, Annals Phys. 326, 2737
(2011); L. Abreu, D. Cabrera and J. M. Torres-Rincon, arXiv:1211.1331 [hep-ph].

[15] H. van Hees and R. Rapp, Nucl. Phys. A 806 (2008) 339

[16] M. Gonin et al. [NA50 Collaboration], Nucl. Phys. A 610, 404C (1996); M. C. Abreu et al.
[NA50 Collaboration], Phys. Lett. B 410, 337 (1997);B. Alessandro et al. [NA50 Collaboration],
Eur. Phys. J. C 39, 335 (2005).

[17] R. Arnaldi et al. [NA60 Collaboration], Nucl. Phys. A 774, 711 (2006), Phys. Rev. Lett. 99,
132302 (2007).

[18] A. Adare et al. [PHENIX Collaboration], Phys. Rev. Lett. 98, 172301 (2007).

[19] T. Matsui and H. Satz, Phys. Lett. B 178, 416 (1986).
[20] J. P. Blaizot and J. Y. Ollitrault, Phys. Rev. Lett. 77, 1703 (1996).

[21] K.L. Haglin, Phys. Rev. C 61 (2000) 031902.

[22] W. Liu, Ph.D. Thesis, (2004) AAT 3157454 (ISBN: 049689885X) http://inspirehep.net/record/707716 Z. Lin and C. M. Ko Phys. Rev. C 62, 034903 (2000).

[23] R. Molina, C. W. Xiao, and E. Oset, Phys. Rev. C 86, 014604 (2012)

[24] B. Friman, S. H. Lee, T. Song, Phys. Lett. B 548 (2002) 153.

[25] A. Hayashigaki, Phys. Lett. B 487, 96 (2000).

[26] T. Hilger, R. Thomas, and B. Kampfer, Phys. Rev. C 79, 025202 (2009); T. Hilger, B. Kampfer and S. Leupold, Phys. Rev. C 84 (2011) 045202

[27] P. Morath, W. Weise, and S. H. Lee, in Proceedings of the 17th Autumn school on QCD: Perturbative or Nonperturbative, Lisbon 1999, edited by L. S. Ferreira, P. Nogueira, and J. I. Silva-Marcos (World Scientific, Singapore, 2001), p. 425.

[28] K. Tsushima, D. H. Lu, A. W. Thomas, K. Saito, and R. H. Landau, Phys. Rev. C 59, 2824 (1999); A. Sibirtsev, K.Tsushima, and A. W. Thomas, Eur. Phys. J. A 6, 351 (1999).

[29] F. O. Gottfried and S. P. Klevansky, Phys. Lett. B 286, 221 (1992).

[30] D. Blaschke, P. Costa, Yu.L. Kalinovsky Phys. Rev. D 85, 034005 (2012); D. Blaschke, G. Burau, Yu. L. Kalinovsky and V. L. Yu-dichev, Prog. Theor. Phys. Suppl. 149, 182 (2003).

[31] L. Tolos, A. Ramos and T. Mizutani, Phys. Rev. C 77, 015207 (2008); L. Tolos, C. Garcia-Recio and J. Nieves, Phys. Rev. C 80, 065202 (2009); L. Tolos, J. Schaffner-Bielich, H. Stöcker, Phys. Lett. B 635 (2006) 85.

[32] A. Sibirtsev, K. Tsushima and A.W. Thomas, Phys.Lett. B 484 (2000) 23.

[33] L. Tolos, J. Schaffner-Bielich, A. Mishra Phys. Rev. C 70, 025203 (2004); Eur. Phys. J. C 43 (2005) 127; J. Phys. G 31 (2005) S1213

[34] G. R. G. Burau, D. B. Blaschke and Y. L. Kalinovsky, Phys. Lett. B 506, 297 (2001)

[35] M.F.M. Lutz and C.L. Korpa, Phys. Lett. B 633, 43 (2006).

[36] T. Mizutani, A. Ramos, Phys.Rev. C 74 (2006) 065201.

[37] R. Molina, D. Gamermann, E. Oset, and L. Tolos, Eur. Phys. J. A 42, 31 (2009)

[38] L.S. Geng, N. Kaiser, J.Martin-Camalich and W.Weise, Phys. Rev. D 82, 054022 (2010)

[39] H.A. Weldon Phys. Rev.D 28, 8 (1983).

[40] S. Ghosh, S. Sarkar and S. Mallik, Eur. Phys. J. C 70, 251 (2010)

[41] M. Le Bellac, Thermal Field Theory (Cambridge University Press, Cambridge, 1996).

[42] R. L. Kobes and G. W. Semenoff, Nucl. Phys. B 260 (1985) 714.

[43] S. Mallik and S. Sarkar, Eur. Phys. J. C 61 (2009) 489
[44] S. Ghosh and S. Sarkar, Nucl. Phys. A 870, 94 (2011).

[45] C. Fuchs, B.V. Martemyanov, A. Faessler and M. I. Krivoruchenko, Phys. Rev. C 73, 035204 (2006).

[46] G. Krein, J. Phys. Conf. Ser. 422 (2013) 012012; K. Tsushima, D.H. Lu, G. Krein, A.W. Thomas, Phys.Rev. C83 (2011) 065208.