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$X(3872)$ production and absorption in a hot hadron gas

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ABSTRACT

We calculate the time evolution of the $X(3872)$ abundance in the hot hadron gas produced in the late stage of heavy ion collisions. We use effective field Lagrangians to obtain the production and dissociation cross sections of $X(3872)$. In this evaluation we include diagrams involving the anomalous couplings $\pi D^* \bar{D}^*$ and $X \bar{D}^* D^*$ and also the couplings of the $X(3872)$ with charged D and D^* mesons. With these new terms the $X(3872)$ interaction cross sections are much larger than those found in previous works. Using these cross sections as input in rate equations, we conclude that during the expansion and cooling of the hadronic gas, the number of $X(3872)$, originally produced at the end of the mixed QGP/hadron gas phase, is reduced by a factor of 4.

1. Introduction

Over the last decade dozens of new charmonium states have been observed [1–4]. Among these new states, the most studied one is the $X(3872)$. It was first observed in 2003 by the Belle Collaboration [5,6], and has been confirmed by other five experiments: BaBar [7], CDF [8–10], DØ [11], LHCb [12,13] and CMS [14]. The LHCb Collaboration has determined the $X(3872)$ quantum numbers to be $J^{PC} = 1^{++}$, with more than 8σ significance [13].

The structure of the new charmonium states has been subject of an intense debate. In the case of $X(3872)$, calculations using constituent quark models give masses for possible charmonium states with $J^{PC} = 1^{++}$ quantum numbers, which are much bigger than the observed $X(3872)$ mass: $2\,^3P_1(3990)$ and $3\,^3P_1(4290)$ [15]. These results, together with the observed isospin violation in their hadronic decays, motivated the conjecture that these objects are multiquark states, such as mesonic molecules or tetraquarks. Indeed, the vicinity of the X mass to the $\bar{D}D^*$ threshold inspired the proposal that the $X(3872)$ could be a molecular $\bar{D}D^*$ bound state with a small binding energy [16,17]. Another well known interpretation of the $X(3872)$ is that it can be a tetraquark state resulting from the binding of a diquark and an antidiquark [18,19]. There are other proposals as well [2–4]. One successful approach in describing experimental data is to treat the X as an admixture of two and four-quark states [20].

Until now it has not been possible to determine the structure of the X, since the existing data on masses and decay widths can be explained by quite different models. However this situation can change, as we address the production of exotic charmonium in hadronic reactions, i.e. proton–proton, proton–nucleus and nucleus–nucleus collisions. Hadronic collisions seem to be a promising testing ground for ideas about the structure of the new states. In Refs. [3,21] it has been argued that it is extremely difficult to produce meson molecules in p p collisions. In these papers the estimated cross section for $X(3872)$ production is two orders of magnitude smaller than the measured one. On the other hand, in Ref. [22] the authors come to a very different conclusion. The subject is still under debate. In Ref. [23] a simple model was proposed to compute the X production cross section in p p collisions in the tetraquark approach. Predictions were made for the next run of the LHC.

As pointed out in Refs. [24,25], high energy heavy ion collisions offer an interesting scenario to study the production of multiquark states. In these processes, a quite large number of heavy quarks are expected to be produced, reaching as much as 20 $c\bar{c}$ pairs per unit rapidity in Pb + Pb collisions at the LHC. Moreover, the formation of quark gluon plasma (QGP) may enhance the production of exotic charmonium states, since the charm quarks are free to move over a large volume and they may coalesce to form bound states at the

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end of the QGP phase or, more precisely, at the end of the mixed phase, since the QGP needs some time to hadronize.

One of the main conclusions of Refs. [24,25] was that, if the production mechanism is coalescence, then the production yield of these exotic hadrons at the moment of their formation strongly reflects their internal structure. In particular it was shown that in the coalescence model the production yield of the $X(3872)$, at RHIC or LHC energies, is almost 20 times bigger for a molecular structure than for a tetraquark configuration.

After being produced at the end of the quark gluon plasma phase, the $X(3872)$ interacts with other hadrons during the expansion of the hadronic matter. Therefore, the $X(3872)$ can be destroyed in collisions with the comoving light mesons, such as $X + \pi \to \bar{D} + D$, $X + \pi \to \bar{D}^* + D^*$, but it can also be produced through the inverse reactions, such as $D + \bar{D} \to X + \pi$, $\bar{D}^* + D^* \to \pi + X$. We expect these cross sections to depend on the spatial configuration of the $X(3872)$. Charm tetraquarks in a diquark–antidiquark configuration $(cq) - (\bar{c}\bar{q})$ have a typical radius comparable to (or smaller than) the radius of the charmonium groundstates, i.e. $r_{4q} \simeq r_{\bar{c}c} \simeq 0.3$–$0.4$ fm. Charm meson molecules are bound by meson exchange and hence $r_{mol} \simeq 1/m_\pi \simeq 1.5$ fm. In fact, the calculations of Ref. [26] show that $r_{mol} \simeq 2.0$–3.0 fm. Along a different line of argumentation and making use of the smallness of the X binding energy, in Ref. [28] the authors conclude that the X radius could be as large as $4.9^{+13.4}_{-1.4}$ fm.

Molecules are thus much bigger than tetraquarks and their absorption cross sections may also be much bigger. On the other hand, when these states are produced from $D - \bar{D}^*$ fusion in a hadron gas, what matters is the overlap between the initial and final state configurations. Assuming that the radius of the D and D^* mesons is $r_D \simeq 0.6$ fm [29], the initial $D + D^*$ state has a larger spatial overlap with a molecule than with a tetraquark and, therefore, the production of molecules is favored. Hence from geometrical arguments we expect that in a hot hadronic environment molecules are easier to produce and also easier to destroy than tetraquarks. Of course geometrical estimates of cross sections are more reliable if we apply them to high energy collisions. Here the typical collision energies are of the order of the temperature $T \simeq 100$–180 MeV and are probably not high enough. Nevertheless, at a qualitative level, they can be useful as guidance.

In Ref. [30] the interactions of the X in a hadronic medium were studied in the framework of $SU(4)$ effective Lagrangians. The authors computed the corresponding production and absorption cross sections, finding that the absorption cross section is two orders of magnitude larger than the production one. The effective Lagrangians include the X particle as a fundamental degree of freedom and the theory is unable to distinguish between molecular and tetraquark configurations. Presumably this information might be included in the form factors introduced in the vertices. The authors find that it is much easier to destroy the X than to create it. In particular, for the largest thermally averaged cross sections they find: $< \sigma v >_{\pi X \to D^* \bar{D}^*} \simeq 30 < \sigma v >_{D^* \bar{D}^* \to \pi X}$. In spite of this difference, the authors of Ref. [30] arrived at the intriguing conclusion that the number of X's stays approximately constant during the hadronic phase. In Ref. [30] the coupling of the $X(3872)$ with charged charm mesons (such as $D^- D^{*+}$) was neglected and only neutral mesons were considered (such as $D^0 \bar{D}^{0*}$). Moreover, the terms with anomalous couplings were not included in the calculations. In Ref. [31] we showed that the inclusion of the couplings of the $X(3872)$ to charged D's and D^*'s and those of the anomalous vertices, $\pi D^* \bar{D}^*$ and $X D^* \bar{D}^*$, increases the cross sections by more than one order of magnitude. Similar results were also observed in the case of $J/\psi \pi$ cross section [32]. These anomalous vertices also give rise to new reaction channels, namely, $\bar{D} + D^* \to \pi + X$ and $\pi + X \to \bar{D} + D^*$. Thus it is important to evaluate the changes that

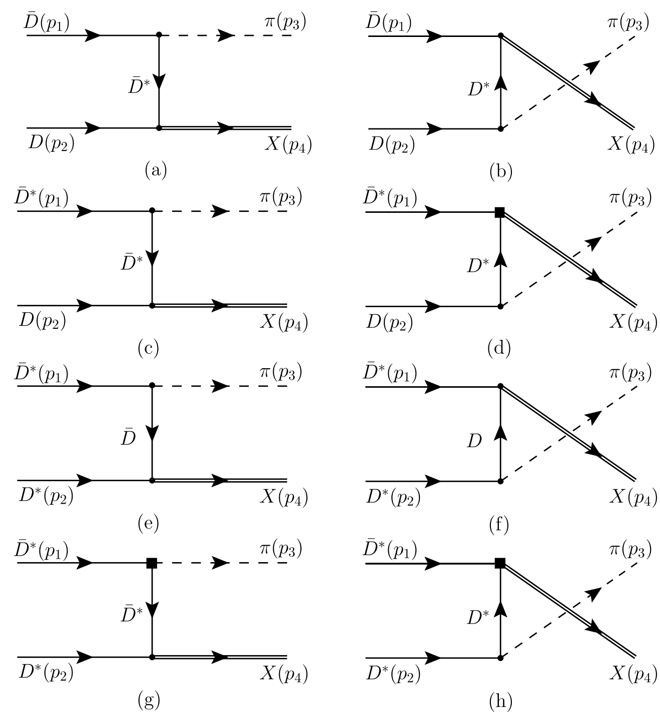

Fig. 1. Diagrams contributing to the processes $\bar{D}D \to \pi X$ [(a) and (b)], $\bar{D}^* D \to \pi X$ [(c) and (d)] and $\bar{D}^* D^* \to \pi X$ [(e), (f), (g) and (h)]. The filled box in the diagrams (d), (g) and (h) represents the anomalous vertex $X D^* \bar{D}^*$, which was evaluated in Ref. [31]. The charges of the particles are not specified.

the above mentioned contributions can produce in the X abundance (and in its time evolution) in reactions as those considered in Ref. [30]. This is the subject of this work.

The formalism used in Refs. [30] and [31] was originally developed to study the interaction of charmonium states (specially the J/ψ) with light mesons in a hot hadron gas many years ago [32, 33]. The conclusions obtained in the past can help us now, giving some baseline for comparison. For example, if it is true that the X has a large $c\bar{c}$ component, we may expect that the corresponding production and absorption cross sections are comparable to the J/ψ ones. If, alternatively, they turn out to be much larger, this could be an indication of a strong multiquark and possibly molecular component.

The paper is organized as follows. In the next section we discuss the cross sections averaged over the thermal distributions. In Sec. 3 we investigate the time evolution of the $X(3872)$ abundance by solving the kinetic equation based on the phenomenological model of Ref. [30]. Finally in Sec. 4 we present our conclusions.

2. Cross sections averaged over the thermal distribution

In this section we calculate the cross sections averaged over the thermal distributions for the processes $\bar{D}D \to \pi X$, $\bar{D}^* D \to \pi X$ and $\bar{D}^* D^* \to \pi X$, and for the inverse reactions. This information is the input to the study of the time evolution of the $X(3872)$ abundance in hot hadronic matter. In Fig. 1 we show the different diagrams contributing to each process. In Ref. [30] it was shown that the contribution from the reactions involving the ρ meson is very small compared to the reactions with pions and thus we neglect the former in what follows. The cross sections for the processes shown in Fig. 1 were obtained in Ref. [31]. To calculate the amplitudes of the processes shown in Fig. 1 we need Lagrangians to determine the contribution of the Pseudoscalar–Pseudoscalar–Vector (PPV), Vector–Vector–Pseudoscalar (VVP) and

Vector–Vector–Vector (VVV) vertices. This can be done considering Lagrangians built using an effective theory in which the vector mesons are identified as the dynamical gauge bosons of the hidden $U(3)_V$ local symmetry in the $U(3)_L \times U(3)_R/U(3)_V$ non-linear sigma model [34–37]. They read:

$$\mathcal{L}_{PPV} = -ig_{PPV} \langle V^\mu [P, \partial_\mu P] \rangle, \tag{1}$$

$$\mathcal{L}_{VVP} = \frac{g_{VVP}}{\sqrt{2}} \epsilon^{\mu\nu\alpha\beta} \langle \partial_\mu V_\nu \partial_\alpha V_\beta P \rangle \tag{2}$$

$$\mathcal{L}_{VVV} = ig_{VVV} \langle (V^\mu \partial_\nu V_\mu - \partial_\nu V_\mu V^\mu) V^\nu \rangle. \tag{3}$$

The Lagrangians above can be extended to SU(4) considering P and V_μ as matrices containing the 15-plet of pseudoscalar and vector mesons and the singlet of SU(4), which in the physical basis and considering ideal mixing for η and η' as well as for ω and ϕ read as [38]:

$$P = \begin{pmatrix} \frac{\eta}{\sqrt{3}} + \frac{\eta'}{\sqrt{6}} + \frac{\pi^0}{\sqrt{2}} & \pi^+ & K^+ & \bar{D}^0 \\ \pi^- & \frac{\eta}{\sqrt{3}} + \frac{\eta'}{\sqrt{6}} - \frac{\pi^0}{\sqrt{2}} & K^0 & D^- \\ K^- & \bar{K}^0 & -\frac{\eta}{\sqrt{3}} + \sqrt{\frac{2}{3}}\eta' & D_s^- \\ D^0 & D^+ & D_s^+ & \eta_c \end{pmatrix}, \tag{4}$$

$$V_\mu = \begin{pmatrix} \frac{\omega+\rho^0}{\sqrt{2}} & \rho^+ & K^{*+} & \bar{D}^{*0} \\ \rho^- & \frac{\omega-\rho^0}{\sqrt{2}} & K^{*0} & D^{*-} \\ K^{*-} & \bar{K}^{*0} & \phi & D_s^{*-} \\ D^{*0} & D^{*+} & D_s^{*+} & J/\psi \end{pmatrix}_\mu . \tag{5}$$

The $\mathcal{L}_{VVP}$ Lagrangian written above describes an anomalous vertex, which involves a violation of the natural parity. The natural parity of a particle is defined for bosons only and it is $P_n = P(-1)^J$, where P is the intrinsic parity and J is the spin of the particle. In other words, the natural parity of a particle is $+1$ if the particle transforms as a true Lorentz-tensor of that rank, and -1 if it transforms as a pseudotensor. In this way the field V has natural parity $+1$, since it represents a vector, but the field P has natural parity -1, since it corresponds to a pseudoscalar. There exists a unique way to construct the interaction Lagrangian that would violate the natural parity and would simultaneously conserve the intrinsic parity and would be Lorentz invariant: by using the Levi-Civita pseudotensor. Therefore anomalous processes are described by the following Lagrangian densities [39,40]:

$$\mathcal{L}_{XD^*D^*} = i\, g_{XD^*D^*}\, \epsilon^{\mu\nu\alpha\beta}\, \partial_\mu X_\nu\, \bar{D}^*_\alpha\, D^*_\beta$$

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

In SU(3), the couplings appearing in Eqs. (1), (2) and (3) are given by [41–43]

$$g_{PPV} = \frac{m_V}{2f_\pi},$$

$$g_{VVP} = \frac{3m_V^2}{16\pi^2 f_\pi^3},$$

$$g_{VVV} = \frac{m_V}{2f_\pi},$$

with m_V being the mass of the vector meson, which we take as the mass of the ρ meson, and $f_\pi = 93$ MeV is the pion decay constant. The symbol $\langle\rangle$ in Eqs. (1), (2) and (3) indicates the trace in the isospin space.

The SU(4) symmetry is not a good symmetry in quantum chromodynamics, since the charm quark is much heavier than the u, d and s quarks. However it turns out that the SU(4) symmetry relations for couplings constants are still meaningful, as shown in QCD sum rules calculations [44]. The main idea of using the SU(4) symmetry here is to classify all the possible interaction vertices among the meson multiplets and then estimate their respective couplings, trying to restrict them as much as possible by using available experimental information. For instance, the g_{PPV} coupling for the $D^*D\pi$ vertex is corrected by considering the heavy quark spin symmetry following [45] to

$$g_{PPV} = \frac{m_V}{2f\pi} \frac{m_{D^*}}{m_{K^*}}. \tag{6}$$

Using effective Lagrangians based on SU(4), the coupling of the $X(3872)$ to $\bar{D}^*D^*$ was estimated through the evaluation of triangular loops and an effective Lagrangian was proposed to describe this vertex. As a result, it was found that the contributions coming from the coupling of the $X(3872)$ to charged D's and D^*'s and from the anomalous vertices play an important role in determining the cross sections. The coupling constant of the $X\bar{D}^*D^*$ vertex was found to be $g_{X\bar{D}^*D^*} = 12.5 \pm 1.0$. For more details about the calculations, we refer the reader to Ref. [31]. In the present manuscript, we follow Ref. [31] and obtain the cross sections of the processes in Fig. 1 using a form factor of the type

$$F(\vec{q}) = \frac{\Lambda^2}{\Lambda^2 + \vec{q}^2}, \tag{7}$$

where $\Lambda = 2.0$ GeV is the cutoff and $\vec{q}$ the three-momentum transfer in the center of mass frame. Following Refs. [30,46], the thermally averaged cross section for a process $ab \to cd$ can be calculated using the expression

$$\begin{aligned} \langle \sigma_{ab\to cd}\, v_{ab} \rangle &= \frac{\int d^3\mathbf{p}_a d^3\mathbf{p}_b\, f_a(\mathbf{p}_a)\, f_b(\mathbf{p}_b)\, \sigma_{ab\to cd}\, v_{ab}}{\int d^3\mathbf{p}_a d^3\mathbf{p}_b\, f_a(\mathbf{p}_a)\, f_b(\mathbf{p}_b)} \\ &= \frac{1}{4\alpha_a^2 K_2(\alpha_a)\alpha_b^2 K_2(\alpha_b)} \int_{z_0}^{\infty} dz K_1(z)\, \sigma(s = z^2T^2) \\ &\quad \times \left[z^2 - (\alpha_a + \alpha_b)^2\right]\left[z^2 - (\alpha_a - \alpha_b)^2\right], \end{aligned} \tag{8}$$

where f_a and f_b are Bose–Einstein distributions, $\sigma_{ab\to cd}$ are the cross sections computed in [31], v_{ab} represents the relative velocity of the two interacting particles (a and b), $\alpha_i = m_i/T$, with m_i being the mass of particle i and T the temperature, $z_0 = max(\alpha_a + \alpha_b, \alpha_c + \alpha_d)$, and K_1 and K_2 are the modified Bessel functions of first and second kind, respectively. The masses used in the present work are: $m_D = 1867.2$ MeV, $m_{D^*} = 2008.6$ MeV, $m_\pi = 137.3$ MeV and $m_X = 3871.7$ MeV [1].

In Fig. 2a we show the thermally averaged cross section for the process $\bar{D}D \to \pi X(3872)$, considering only the coupling of the $X(3872)$ to the neutral states $\bar{D}^0 D^{*0}$ (solid line) and adding the coupling to the charged components (dashed line). As can be seen, the thermally averaged cross section increases by a factor of about 2.5 when we include the charged D's and D^*'s.

In Figs. 2b and 2c we show the thermally averaged cross sections for the processes $\bar{D}^*D \to \pi X(3872)$ and $\bar{D}^*D^* \to \pi X(3872)$, considering only the coupling of the X to neutral D's and D^*'s (dashed line), including couplings to charged D's and D^*'s (dotted line) and finally adding also the contribution from the anomalous vertices (shaded region). As can be seen, the contribution from the anomalous vertices produces an enhancement of the thermally averaged cross sections by a factor of 100–150.

In Fig. 3a we show the total thermally averaged cross sections for the processes involving the production of the $X(3872)$ state, i.e., $\bar{D}D \to \pi X$, $\bar{D}^*D \to \pi X$ and $\bar{D}^*D^* \to \pi X$ reactions,

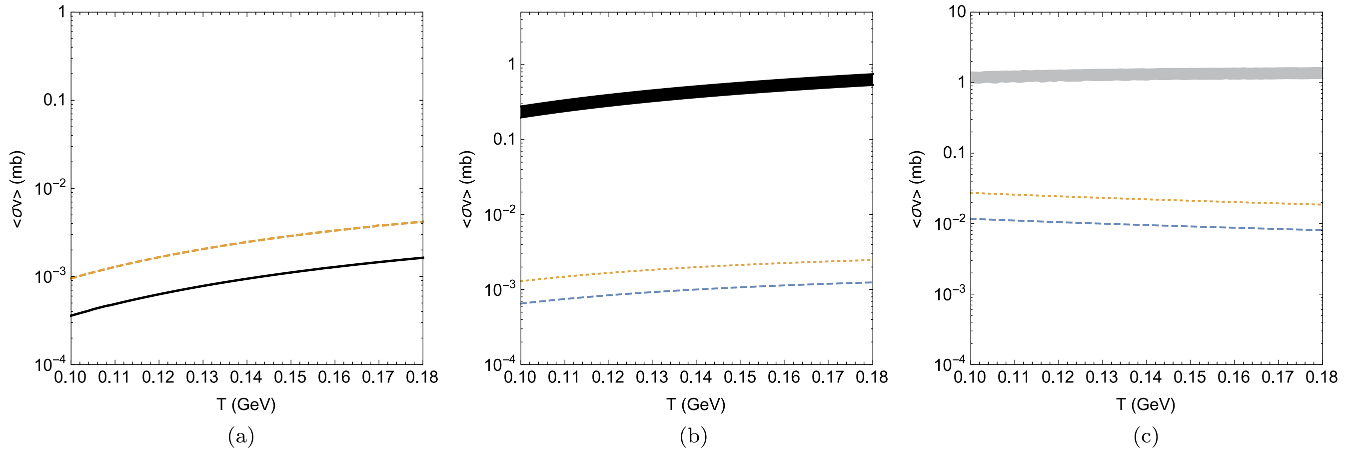

Fig. 2. Thermally averaged cross sections. a) $\bar{D}D \to \pi X(3872)$, considering only the coupling of the X to the neutral D's and D^*'s (solid line) and adding the coupling to charged D's and D^*'s (dashed line). b) $\bar{D}^*D \to \pi X(3872)$, considering only the coupling of the X to neutral D's and D^*'s (dashed line), including the contribution from charged D's and D^*'s (dotted line) and also including diagrams containing the anomalous vertices (shaded region). c) $\bar{D}^*D^* \to \pi X(3872)$. The lines and shaded area have the same meaning as those of b).

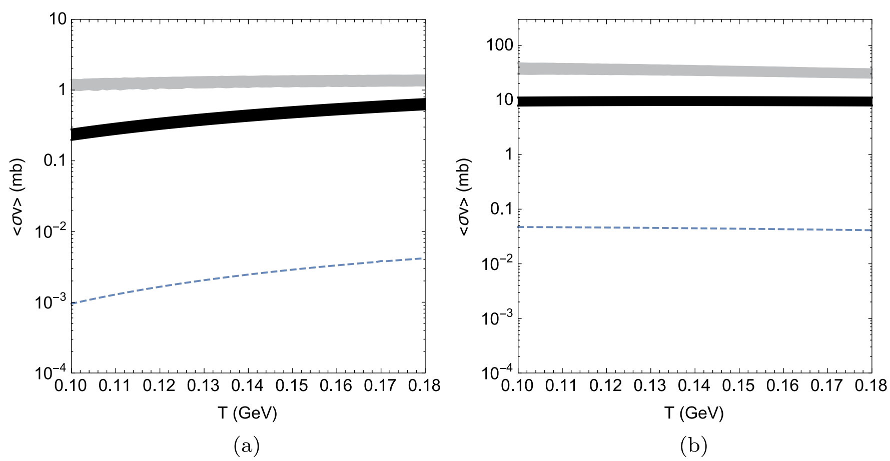

Fig. 3. Thermally averaged cross sections. a) $\bar{D}D \to \pi X(3872)$ (dashed line), $\bar{D}^*D \to \pi X(3872)$ (dark-shaded region) and $\bar{D}^*D^* \to \pi X(3872)$ (light-shaded region). b) $\pi X(3872) \to \bar{D}D$ (dashed line), $\pi X(3872) \to \bar{D}^*D$ (dark-shaded region) and $\pi X(3872) \to \bar{D}^*D^*$ (light-shaded region).

while in Fig. 3b we show the inverse processes, i.e., the dissociation of $X(3872)$ through the reactions $\pi X \to \bar{D}D$, $\pi X \to \bar{D}^*D$, $\pi X \to \bar{D}^*D^*$, respectively. For the latter cases, we use the principle of detailed balance to determine the corresponding cross sections. Fig. 3 should be compared with the Fig. 3 of Ref. [30]. Our cross sections are a factor 100 larger in reactions with $D^*\bar{D}^*$ in the initial or final state This can be attributed mostly to the inclusion of the anomalous terms. In [31] it was found that $g_{XD^*D^*} \simeq 12$. This number, when squared in the calculation of the cross sections, yields a factor larger than 100. A similar change due to the contribution of anomalous vertices was seen in the past, in the study of the J/ψ interaction with light hadrons [27]. Our cross sections are a factor 10 larger in the case of DD mesons in the initial or final state. The difference comes from the inclusion of the coupling of the X to charged D's and D^*'s, which was not considered in Ref. [30]. On the other hand, in both works, the absorption cross sections are more than fifty times larger than the production ones.

Our results can also be compared with those obtained in [28], which describes an interesting attempt to study the reactions $D^*\bar{D}^* \to \pi X(3872)$ and $\pi X(3872) \to D^*\bar{D}^*$ with an improved version of the XEFT (X Effective Field Theory). In [28] the formalism employed relies on low energy approximations and it is applicable when the kinetic collision energy is much smaller than the pion mass. In our approach we can use the theory at any collision energy. The collision energy relevant for the X interactions in heavy ion collisions is of the order of the fireball temperature, i.e., $\simeq$ 100–170 MeV. Moreover, the reactions $D\bar{D} \to \pi X(3872)$ and specially $D\bar{D}^* \to \pi X(3872)$ (and the corresponding inverse reactions) are extremely important for the computation of the X yield. Unfortunately, as mentioned in [28], these reactions can not be described by the XEFT, because they involve the exchange of a charm meson, which is off its energy shell by an amount of order m_π. In [28] the X is assumed to be a molecular bound state with a small binding energy E_X, whereas in our effective Lagrangian approach the X is an independent degree of freedom and information on its possible molecular nature is hidden in coupling constants, such as g_{XD^*D}. As in [28], in our calculations we also find an enhancement in the $D^*\bar{D}^* \to \pi X(3872)$ cross section at low collision energies. In order to make a comparison with the results found in [28], we have calculated the cross section of the particular process $D^{*+}\bar{D}^{*0} \to X\pi^+$ (where there is an exchange of a D meson), using the central values of the D^{*0}, D^+, X and π^+ physical masses. The values of the coupling constant $g_{D^*D\pi}$ is the one usually found in the literature (see [31]) and $g_{XD^*D} = 3651/\sqrt{2}$ MeV. Changing the values of the X mass, we indirectly change the value of the binding energy. We obtain results which are very close to the ones found in [28].

In the computation of the time evolution of the X abundance, we will need to know how the temperature changes with time and this is highly model dependent. Fortunately, as one can see in Fig. 3, the dependence of $< \sigma v >$ on the temperature is relatively weak.

3. Time evolution of the $X(3872)$ abundance

Following Ref. [30] we study the yield of $X(3872)$ in central Au–Au collisions at $\sqrt{s_{NN}} = 200$ GeV. By using the thermally averaged cross sections obtained in the previous section, we can now analyze the time evolution of the $X(3872)$ abundance in hadronic matter, which depends on the densities and abundances of the particles involved in the processes of Fig. 1, as well as the cross sections associated with these reactions (and the corresponding inverse reactions), Figs. 3a and 3b. The momentum-integrated evolution equation has the form [30,46–48]

$$\frac{dN_X(\tau)}{d\tau} = R_{QGP}(\tau) + \sum_{c,c'}[\langle\sigma_{cc'\to\pi X}v_{cc'}\rangle n_c(\tau)N_{c'}(\tau) - \langle\sigma_{\pi X\to cc'}v_{\pi X}\rangle n_\pi(\tau)N_X(\tau)], \quad (9)$$

where $N_X(\tau)$, $N_{c'}(\tau)$, $n_c(\tau)$ and $n_\pi(\tau)$ are the abundances of $X(3872)$, of charmed mesons of type c', of charmed mesons of type c and of pions at proper time τ, respectively. The term $R_{QGP}(\tau)$ in Eq. (9) represents the X production from the quark–gluon plasma in the mixed phase, since the hadronization of the QGP takes a finite time, and it is given by [30,48]:

$$R_{QGP}(\tau) = \begin{cases} \dfrac{N_X^0}{\tau_H - \tau_C}, & \tau_C < \tau < \tau_H, \\ 0, & \text{otherwise}, \end{cases} \quad (10)$$

where the times $\tau_C = 5.0$ fm/c and $\tau_H = 7.5$ fm/c determine the beginning and the end of the mixed phase respectively. The constant N_X^0 corresponds to the total number of $X(3872)$ produced from quark–gluon plasma. To solve Eq. (9) we assume that the pions and charmed mesons in the reactions contributing to the abundance of X are in equilibrium. Therefore $N_{c'}(\tau)$, $n_c(\tau)$ and $n_\pi(\tau)$ can be written as [30,46–48]

$$\begin{aligned} N_{c'}(\tau) &\approx \frac{1}{2\pi^2}\gamma_C\, g_D\, m_{D^{(*)}}^2\, T(\tau)\, V(\tau)\, K_2\left(\frac{m_{D^{(*)}}}{T(\tau)}\right), \\ n_c(\tau) &\approx \frac{1}{2\pi^2}\gamma_C\, g_D\, m_{D^{(*)}}^2\, T(\tau)\, K_2\left(\frac{m_{D^{(*)}}}{T(\tau)}\right), \\ n_\pi(\tau) &\approx \frac{1}{2\pi^2}\gamma_\pi\, g_\pi\, m_\pi^2\, T(\tau)\, K_2\left(\frac{m_\pi}{T(\tau)}\right), \end{aligned} \quad (11)$$

where γ_i and g_i are the fugacity factor and the spin degeneracy of particle i respectively. As can be seen in Eq. (11), the time dependence in Eq. (9) enters through the parametrization of the temperature $T(\tau)$ and volume $V(\tau)$ profiles suitable to describe the dynamics of the hot hadron gas after the end of the quark–gluon plasma phase. Following Refs. [30,47,48], we assume the τ dependence of T and V to be given by [30,47,48]

$$\begin{aligned} T(\tau) &= T_C - (T_H - T_F)\left(\frac{\tau - \tau_H}{\tau_F - \tau_H}\right)^{\frac{4}{5}}, \\ V(\tau) &= \pi\left[R_C + v_C(\tau - \tau_C) + \frac{a_C}{2}(\tau - \tau_C)^2\right]^2 \tau_C. \end{aligned} \quad (12)$$

These expressions are based on the boost invariant picture of Bjorken [49] with an accelerated transverse expansion. In the above equation $R_C = 8.0$ fm denotes the final size of the quark–gluon plasma, while $v_C = 0.4\,c$ and $a_C = 0.02\,c^2$/fm are its transverse flow velocity and transverse acceleration at τ_C respectively. The critical temperature of the quark gluon plasma to hadronic matter transition is $T_C = 175$ MeV; $T_H = T_C = 175$ MeV is the temperature of the hadronic matter at the end of the mixed phase. The freeze-out takes place at the freeze-out time $\tau_F = 17.3$ fm/c, when the temperature drops to $T_F = 125$ MeV.

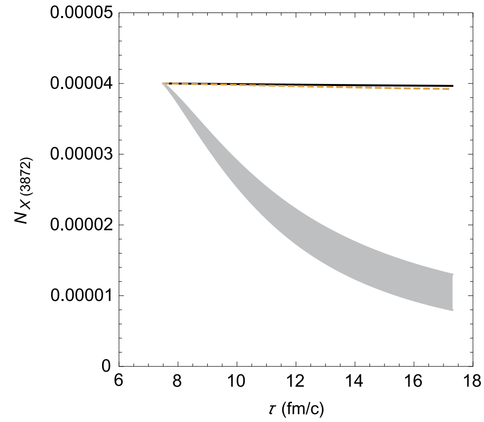

Fig. 4. Time evolution of the $X(3872)$ abundance as a function of the proper time τ in central Au–Au collisions at $\sqrt{s_{NN}} = 200$ GeV. The solid line, the dashed line and the light-shaded region represent the results obtained considering only the neutral D's and D^*'s, adding the contribution from charged D's and D^*'s and including contributions from the anomalous vertices respectively. The initial condition is the abundance of the $X(3872)$ considered as a tetraquark Eq. (13).

To solve Eq. (9), we assume that the total number of charm quarks in charm hadrons is conserved during the production and dissociation reactions, and that the total number of charm quark pairs produced at the initial stage of the collisions at RHIC is 3, yielding the charm quark fugacity factor $\gamma_C \approx 6.4$ in Eq. (11) [25,30]. In the case of pions, we follow Ref. [48] and work with the assumption that their total number at freeze-out is 926, which fixes the value of γ_π appearing in Eq. (11) to be ~ 1.725.

In Ref. [30] the authors studied the yields obtained for the $X(3872)$ abundance within two different approaches: the statistical and the coalescence models. In the statistical model, hadrons are produced in thermal and chemical equilibrium. This model does not contain any information related to the internal structure of the $X(3872)$ and, for this reason we do not consider it in this work. In the case of the coalescence model, the determination of the abundance of a certain hadron is based on the overlap of the density matrix of the constituents in an emission source with the Wigner function of the produced particle. This model contains information on the internal structure of the considered hadron, such as angular momentum, multiplicity of quarks, etc. According to Ref. [25], the number of $X(3872)$ produced at the end of the mixed phase, assuming that the $X(3872)$ is a tetraquark state with $J^{PC} = 1^{++}$, is given by:

$$N^0_{X(4q)} = N_{X(4q)}(\tau_H) \approx 4.0 \times 10^{-5}. \quad (13)$$

In order to determine the time evolution of the $X(3872)$ abundance we solve Eq. (9) starting at the end of the mixed phase, i.e. at $\tau_H = 7.5$ fm/c, and assuming that the $X(3872)$ is a tetraquark, formed according to the prescription of the coalescence model. The initial condition is given by Eq. (13). We use this initial abundance to integrate Eq. (9) and we show the result in Fig. 4. In the figure the solid line represents the result obtained using the same approximations as those made in Ref. [30]. Our curve is slightly different from that of Ref. [30] because we did not include the contribution of the ρ mesons, as discussed earlier. The dashed line shows the result when we include the couplings of the $X(3872)$ to charged D's and D^*'s. The light-shaded band shows the results obtained with the further inclusion of the diagrams containing the anomalous vertices. The band reflects the uncertainty in the $X\bar{D}^*D^*$ coupling constant, which is $g_{X\bar{D}^*D^*} = 12.5 \pm 1.0$ [31].

As can be seen, without the inclusion of the anomalous coupling terms, the abundance of X remains basically constant. This is because the magnitude of the thermally averaged cross sections for the X production and absorption reactions obtained within this approximation is so small that the second term in the right hand side of Eq. (9) is completely negligible compared to the first term. When including the coupling of the X to charged D's and D^*'s we basically do not find any important change for the time evolution of the X abundance, since, as can be seen in Fig. 2, the thermally averaged cross sections do not change drastically in this case. On the other hand, the inclusion of the anomalous coupling terms, depicted in Figs. 1c, 1d, 1f, 1g and 1h, modifies the behavior of the $X(3872)$ yield, producing a fast decrease of the X abundance with time. This is the main result of this work. We emphasize that the $X(3872)$ abundance, whose time evolution was studied above, is the only one which comes from the QGP and is what could be observed if the $X(3872)$ is a tetraquark state. However, if the $X(3872)$ is a molecular state, it will be formed by hadron coalescence at the end of the hadronic phase. According to Ref. [30], at this time the average number of X's, considering it as a $D\bar{D}^*$ molecule, is

$$N_{X(mol)} \approx 7.8 \times 10^{-4}\,, \tag{14}$$

which is about 80 times larger than the contribution for a tetraquark state at the end of the hadronic phase (see Fig. 4). We can then conclude that the QGP contribution for the $X(3872)$ production (as a tetraquark state and from quark coalescence) after being suppressed during the hadronic phase, becomes insignificant at the end of the hadronic phase.

4. Concluding remarks

In this work we have studied the time evolution of the $X(3872)$ abundance in heavy ion collisions. If the $X(3872)$ is a tetraquark state it will be produced at the mixed phase by quark coalescence. After being produced at the end of the quark gluon plasma phase, the $X(3872)$ interacts with other comoving hadrons during the expansion of the hadronic matter. Therefore, the $X(3872)$ can be destroyed in collisions with the comoving light mesons, such as $X + \pi \to \bar{D} + D$, $X + \pi \to \bar{D}^* + D^*$ but it can also be produced through the inverse reactions, such as $D + \bar{D} \to X + \pi$, $\bar{D}^* + D^* \to \pi + X$. In this work we have considered the contributions of anomalous vertices, $\pi D^* \bar{D}^*$ and $X\bar{D}^*D^*$, and the contributions from charged D and D^* mesons to the $X(3872)$ production and dissociation cross sections. These vertices, apart from enhancing the cross sections associated with the $\bar{D}^*D^*$ channel, give rise to additional production/absorption mechanisms of X, which are found to be relevant.

The cross sections, averaged over the thermal distribution, have been used to analyze the time evolution of the $X(3872)$ abundance in hadronic matter. We have found that the abundance of a tetraquark X drops from $N_{X(4q)} \approx 4.0 \times 10^{-5}$ at the beginning of the hadronic phase [30] to $N_{X(4q)} \sim 1.0 \times 10^{-5}$ at the end of the hadronic phase.

On the other hand, if the $X(3872)$ is a molecular state it will be produced by hadron coalescence at the end of the hadronic phase. According to Ref. [30], at this time the average number of X's, considering it as a $D\bar{D}^*$ molecule, is $N_{X(mol)} \approx 7.8 \times 10^{-4}$, which is about 80 times larger than $N_{X(4q)}$.

As expected, the results show that the X multiplicity in relativistic ion collisions depends on the structure of $X(3872)$. Our main conclusion is that the contribution from the anomalous vertices play an important role in determining the time evolution of the $X(3872)$ abundance and they lead to strong suppression of this state during the hadronic phase. Therefore, within the uncertainties of our calculation we can say that if the $X(3872)$ were observed in a heavy ion collision it must have been produced at the end of the hadronic phase and, hence, it must be a molecular state.

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