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A potential approach to the X(3872) thermal behaviour

Miguel A. Escobedo,[a,∗] **Néstor Armesto,**[b] **Elena G. Ferreiro**[b] **and Víctor López-Pardo**[b]

[a] *Departament de Física Quàntica i Astrofísica and Institut de Ciències del Cosmos, Universitat de Barcelona,*
Martí i Franquès 1, 08028 Barcelona, Catalonia, Spain

[b] *Instituto Galego de Física de Altas Enerxías (IGFAE), Universidade de Santiago de Compostela, E-15782, Galicia, Spain*

E-mail: miguel.a.escobedo@fqa.ub.edu

We study the potential of X(3872) at finite temperature in the Born-Oppenheimer approximation under the assumption that it is a tetraquark. We argue that, at large number of colors, it is a good approximation to assume that the potential consists in a real part plus a constant imaginary term. The real part is then computed adapting an approach by Rothkopf and Lafferty and using as input lattice QCD determinations of the potential for hybrids. This model allows us to qualitatively estimate at which temperature range the formation of a heavy tetraquark is possible, and to propose a qualitative picture for the dissociation of the state in a medium. Our approach can be applied to other suggested internal structures for the X(3872) and to other exotic states. This work summarizes the results of [1].

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∗Speaker

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1. Introduction

In nature, there exist quarkonium-like particles whose quantum numbers and properties cannot be explained by the simple quark-antiquark model. Among them, we focus on the $X(3872)$ [2], whose internal structure is still a matter of debate. There are two competing models: the tetraquark and the hadronic molecule. On one hand, the tetraquark is a compact bound state of four quarks, in our case two heavy and two light. On the other hand, a hadronic molecule is formed by two heavy-light mesons joint by the strong force analogue of van der Waals interaction.

In order to understand the internal structure of the $X(3872)$, a common approach is to formulate a theoretical model and then check if it is compatible with the observed properties of the state. These properties include, for example, its quantum numbers, spectroscopy, and decay channels. Recently, the $X(3872)$ has been observed in heavy-ion collisions [3]. This opens the possibility to gain information about the structure of the bound state in a different way. That is, studying how the presence of the quark-gluon plasma modifies its behavior.

In this manuscript, we are going to discuss the modification to the potential of the $X(3872)$ at finite temperature. First, let us review briefly the state-of-art of conventional quarkonium in heavy-ion collisions. In recent years, it has been understood that the potential has both a real and an imaginary part [4]. The origin of this imaginary part is the inelastic collision of quarkonium with medium particles. In order to understand the role of the imaginary part in the evolution of the state of quarkonium the formalism of open quantum system was applied (see [5] for a review). The approach has been used to obtain phenomenological predictions that agree with observations [6]. Using this framework, it was found that solving the Schrödinger equation with a complex potential is a good approximation when the binding energy is much smaller than the temperature and when regeneration effects are not considered. This motivates us to find the potential of the $X(3872)$ at finite temperature.

2. Theoretical framework

We are going to use the Born-Oppenheimer approximation. We assume that heavy quarks move non-relativistically around the center-of-mass, with velocity v. The fact that $\Lambda_{QCD} \gg E \sim m_Q v^2$ (where E is the binding energy) implies that the dynamics of light quarks and gluons is much faster than that of the evolution of the bound state. In other words, from the point of view of the heavy quarks, the light particles move very fast. Therefore, the effect of light particles and gluons can be encoded in a potential computed assuming that the heavy quarks are frozen and separated a given distance r. This involves a two-step approximation. First, we compute the potential taking the heavy quarks as static color sources. After this, we solve the Schrödinger equation with that potential.

First, let us discuss the potential at $T = 0$. Ideally, we would like to use a lattice QCD potential for the tetraquark as a starting point. However, this is not yet available since computations with dynamical light quarks are expensive. Instead, we use hybrid data and make the approximation that the tetraquark potential would behave qualitatively similar to the hybrid potential. An hybrid is a bound state of a quark-antiquark pair in the octet representation with a cloud of gluonic degrees of freedom such that the whole state is color neutral. Therefore, it is similar to a tetraquark in which the role of light quarks is played by gluons. Morevover, we assume a single channel approximation,

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where the heavy quarks do not contribute to the spin of the tetraquark. We do not expect that any of these approximations will change the qualitative picture and a quantitative study is out of the scope of this manuscript. The hybrid lattice potential that we are going to use is taken from [7]. It happens that this potential is well fitted by the following formula

$$V(r,0)=\frac{A_{-1}}{r}+A_0+A_2r^2\,. \tag{1}$$

Note that the previous formula is accurate for the distances probed in the lattice computation and, therefore, useful to study whether or not bound state formation is possible. However, eq. (1) is not valid at large distances in which we know, thanks to effective string theory, that the potential rises linearly.

Next, we must discuss how to extend this potential to finite temperature. We are going to study separately the real and the imaginary part of the potential. For the real part of the potential, we use the approach developed in [8]. There, the authors were able to reproduce the static potential of quarkonium at finite temperature using as input the lattice potential at $T=0$. The real part of the potential is obtained from the following convolution

$$\mathrm{Re}\,V(\mathbf{p})=\mathrm{Re}\left(\frac{V_{vac}(\mathbf{p})}{\epsilon(\mathbf{p},m_D)}\right)\,, \tag{2}$$

where ϵ is the medium permittivity in the HTL approximation, m_D is the Debye mass and V_{vac} is the potential at $T=0$. In our case, starting from eq. (1), we obtain the following real part of the potential at finite temperature

$$\mathrm{Re}\,V(r,m_D)=A_{-1}\left(m_D+\frac{e^{-m_Dr}}{r}\right)+A_0+A_2\left[\frac{6}{m_D^2}(1-e^{-m_Dr})-\left(2r^2+\frac{6r}{m_D}\right)e^{-m_Dr}\right]\,. \tag{3}$$

Now, let us discuss the imaginary part of the potential. In the case of conventional quarkonium, it has the following properties. At short distances, the imaginary potential goes like r^2 because the medium sees quarkonium as a small dipole. At long distances, the heavy quarks are not correlated, so the imaginary part of the potential is equal to $-i$ times the decay width of a single heavy quark. Between these two limits, we expect that the imaginary part of the potential is a smoothly increasing function. Indeed, the imaginary potential of quarkonium computed in the HTL approximation fulfills these properties. $\mathrm{Im}V=-C_F\alpha_sT\phi(rm_D)$, where ϕ is a monotonically increasing function between 0 and 1. Knowing this, it is easy to infer how the imaginary part of the potential of a tetraquark would behave. In this case, the heavy quarks are in an octet state. When $r\to 0$, the medium sees the heavy quark pair as a non-relativistic heavy gluon, and the imaginary part of the potential will be $-i/2$ times the decay width of a heavy gluon. At large distances, the two quarks are uncorrelated, similar to the color singlet case. At intermediate distances, we expect that the imaginary part of the potential is a smooth function that interpolates between these two extreme regimes. However, in the large N_c limit, the decay width of a heavy gluon is equal to that of two heavy quarks. Therefore, we can take the imaginary part of the potential to be a constant. Finally, we have to make an educated guest of the size of this constant. We choose the following

$$\Gamma=A_{-1}T+A_2\frac{T}{m_D^2}\,. \tag{4}$$

The rationale for this formula is the following. We expect that the same parameters that are involved in the real part of the potential also appear here, except A_0 that can be absorbed by a redefinition of the heavy quark mass. The rest is fixed by dimensional analysis. Obviously, this is just an educated guest that will allow us to establish the qualitative importance of the effect.

3. Results

3.1 Dissociation temperature

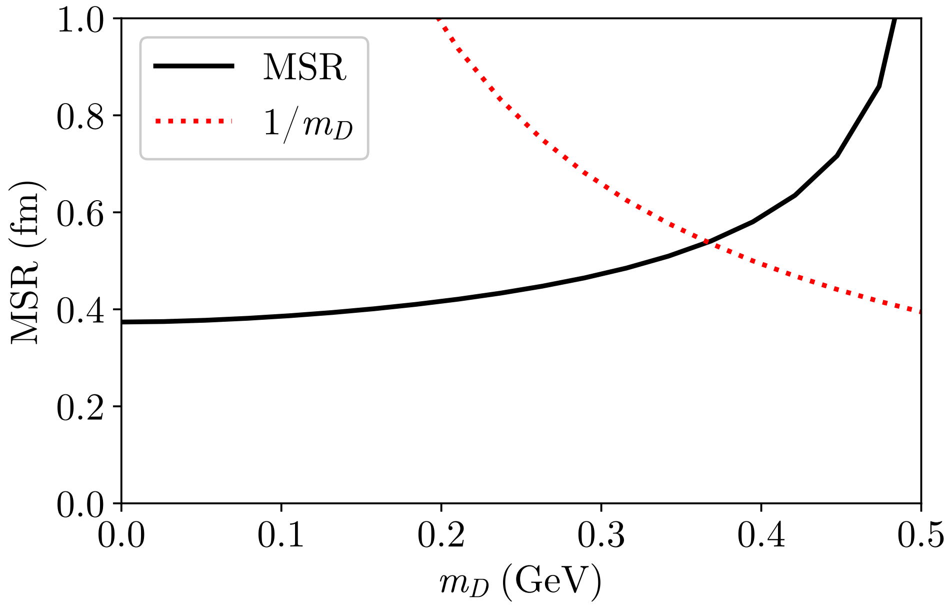

Figure 1: Square root of the mean square radius $\sqrt{\langle r^2 \rangle}$ of the $X(3872)$ as a function of the Debye mass.

The dissociation temperature is the temperature above which the bound state no longer exists in the medium. It is obtained by solving the Schrödinger equation using the complex potential. Since in our case the imaginary part is a constant, it factors out and does not affect the solution. The result that we found is that the dissociation temperature is around $T_d \sim 250\,\text{MeV}$. In fig. 1, we plot $\sqrt{\langle r^2 \rangle}$ against the Debye mass. We observe, that as the temperature approaches T_d (note that $m_D(T_d) \sim 2T_d$) the mean radius of the wave function starts to rise very quickly, signaling the dissapearance of the bound state.

3.2 Survival probability and nuclear modification factor

The survival probability can be computed using the following formula

$$S(t) = \mathrm{e}^{-\int_{t_0}^{t} \Gamma(T(\tau),\tau)\, d\tau} . \tag{5}$$

Note that in our case the survival probability takes a particularly simple form because the decay

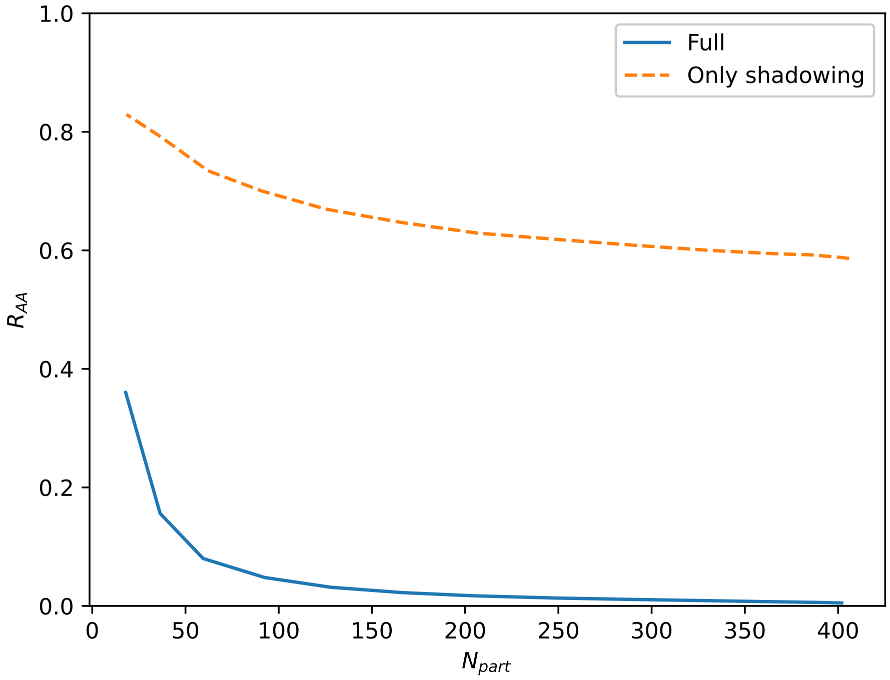

Figure 2: Prediction for R_{AA} of $X(3872)$ at LHCb conditions. The dashed line considers only cold nuclear matter effects, following the model discussed in [9]

width is a constant that does not depend on the form of the wave function. Using this formula, we can compute the nuclear modification factor, R_{AA}, once we know how the temperature seen by the bound state changes with time and position. To do this we follow the lines of [9]. We consider a Bjorken expansion starting at $t = 0.6\,\mathrm{fm}$ and finishing at the time at which the temperature goes below 175 MeV, close to the phase transition. The initial temperature at a given point is computed using a model that takes into account shadowing. If this temperature is larger than T_d then the state is not formed. Otherwise, we compute the survival probability. The results obtained following this procedure can be seen in fig. 2, where we have considered the range of $p_\perp$ and pseudo-rapidity that LHCb can cover for this observale. We observe a strong suppression and only a mild influence of cold nuclear matter effects. We note that, at the moment, we have not yet considered recombination effects.

4. Conclusions

In this manuscript, we have studied exotic quarkonia in heavy-ion collisions with the aim of sheding light on their internal structure. We have developed a qualitative model for the potential of the $X(3872)$. Our results indicate that, for this state, suppression is dominated by screening while only a mild contribution from the imaginary part of the potential is observed.

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