Dynamical Quarkonia Suppression from Complex Potentials

To cite this article: J Casalderrey-Solana 2013 J. Phys.: Conf. Ser. 422 012009

View the article online for updates and enhancements.

Related content
- Heavy flavor and Quarkonia in heavy-ion collisions with the CMS
  Hyunchul Kim and the CMS Collaboration
- Quarkonia production in relativistic heavy ion collisions
  Che Ming Ko, Kyongchol Han and Taesoo Song
- Probing strongly interacting matter with heavy resonances in Pb+Pb collisions at LHC energies
  Prashant Shukla and Abdulla Abdulsalam
Dynamical Quarkonia Suppression from Complex Potentials

J. Casalderrey-Solana
Departament d’Estructura i Constituents de la Matèria and Institut de Ciències del Cosmos (ICCUB), Universitat de Barcelona, Martí i Franquès 1, 08028 Barcelona, Spain
E-mail: jorge.casalderrey@ub.edu

Abstract. I address the effect that a temperature dependent potential has on the suppression of heavy quarkonia states in deconfined hadronic matter. I focus on a simple medium: a homogeneous, fixed temperature and deconfined system with a finite lifetime (QGP-brick). Assuming that all the interactions of a heavy quark anti-quark \((Q-\bar{Q})\) pair with the medium can be recast into an in-medium potential, I solve the time dependent Schrödinger equation to evolve the density matrix which describes the hard pair production and its connection to the final distribution of hadrons after the medium disappears. For those temperatures in which bound states survive in the medium, I find a non-trivial dependence of the production of excited quarkonia states on the in-medium levels, due to the mixing of vacuum and in-medium wave functions. This mixing leads, in particular, to the enhancement of the relative abundance of 2S to 1S states for the case in which the in-medium ground state of the \(Q-\bar{Q}\) system is dissolved.

1. Introduction
The physics of quarkonia suppression in heavy ion collisions has experienced a radical change in the last few years. On the one hand, experimental advances at the LHC (and also at RHIC) have opened access to the bottomonium family. On the theory side, the survival of quarkonia above deconfinement together with the realization that the static properties of in-medium bound states must include complex heavy quark potentials have radically changed the picture of quarkonia dynamics in a deconfined plasma. In this proceedings I will explore the effects that these two features have on the pattern of production of the different quarkonia states.

Following [1], I will study quarkonia production in a homogeneous medium at fixed temperature with a finite lifetime after which the temperature drops abruptly to zero. The spatial extent of the medium is taken to be large, much larger than the in-medium radius of bound states. I will assume that a \(Q-\bar{Q}\) pair is formed within the medium via a hard process at a time \(t_0 = 0\) and interacts with it during a finite time \(t\) after which the medium disappears. After this time, the evolution occurs as in vacuum. In analogy with the jet quenching studies, where a similar set-up was used to analyze the differences among the available quenching models, I will call this set-up, which I have sketched in fig. (1), a QGP-brick.

2. Formalism
From the point of view of the soft dynamics, the hard production of the \(Q-\bar{Q}\) can be described via a Wigner function, \(W_{HI}(r,q)\). Independently of the particular process that leads to the
Figure 1. Sketch of the QGP-Brick set up. At initial time a $Q - \bar{Q}$ pair is formed in the medium via a hard process that occurs within a short time $t_H \sim 1/M$. At a later time of order the inverse binding energy, the soft interactions lead to the formation of quarkonia states from the pair, which are influenced by the presence of the medium. At a later time $t$, the medium disappears and the $Q - \bar{Q}$ system projects into vacuum quarkonia states.

Production of a $Q - \bar{Q}$ pair within the medium, since the process is hard, the typical momentum of each of the quarks, $p$, is large $M < p$. Thus, the typical relative quark momentum is also much larger than the inverse bound state radius and, from the point of view of the soft matrix element, only the $q \to 0$ limit of $W_H$ is relevant for the production of bound states, which after a Fourier transform, leads to an approximate initial density matrix given by

$$\rho_H \left( r - \frac{y}{2}, r + \frac{y}{2}; t_0 \right) = \delta(r)\delta(y),$$

where I have assumed that the hard process takes place at very small distances.

As the $Q - \bar{Q}$ pair propagates, the density matrix eq (1) evolves via a non-relativistic hamiltonian, since the late time interactions are soft as compared to the pair mass. At any given time $t > t_0$ the different particle yields can be obtained by projecting the evolved density matrix into the single state density matrix

$$Y_S \propto \text{Tr} \left( \rho_H(t)\rho_s \right),$$

with $\rho_s = |S\rangle \langle S|$ and $|S\rangle$ the quarkonium estate of interest. Using the cyclicity of the trace, the forward time-evolution of the initial density matrix can be transformed to a backwards evolution of the final quarkonium wave function. Defining $\tilde{t} = t - t_0$ this wave function satisfies the (hermitian conjugate) Schrödinger equation:

$$-i\partial_{\tilde{t}}\tilde{\psi}(\tilde{t}, r) = -\nabla^2 \tilde{\psi}(\tilde{t}, r) + V^\dagger(r)\tilde{\psi}(\tilde{t}, r),$$

with $\tilde{\psi}(\tilde{t}, r) = \langle r|\tilde{S}\rangle$ and initial condition $\tilde{\psi}(0, r) = \langle r|S\rangle$. Combining eq. (1) and eq. (2), the yield of quarkonia production after a time of propagation $t$ is given by

$$Y_S \propto \left| \tilde{\psi}(t - t_0, r = 0) \right|^2,$$

which, in the vacuum, coincides with the relevant long distance matrix element for production in the singlet model. For production in the medium-brick, the interaction with the QGP fields...
changes the potential during the time extent of the medium and changes abruptly back to the vacuum expression when the $Q - \bar{Q}$ leaves the QGP-brick. The complicated dynamics of the production process are given by the time dependent Schrödinger equation, eq. (3). In this study I will focus in a complex potential of the form.

$$V(r) = -\alpha \left( e^{-\mu_D r} \frac{e^{-\mu_D r}}{r} + i\mu_D \phi(\mu_D r) \right), \quad (5)$$

where $\mu_D$ is the Debye screening length of the plasma and

$$\phi(r) = \int_0^\infty dz \frac{z}{(z^2 + 1)^2} \left( 1 - \frac{\sin(zr)}{zr} \right), \quad (6)$$

which coincides with the functional form of the imaginary part of the perturbative potential [2, 3]. To understand the effect of this imaginary part on the yield, I will compare the results of this potential with a real potential given by $\text{Re} V(r)$.

Since after the particle leaves the brick the yield of the different states remains constant, the modification of the quarkonia yields is controlled by the evolution of the hard density matrix up to the time $t$ when the medium disappears. The yield suppression is, then, given by ratio of yields in medium and in vacuum at this time.

$$R_S = \frac{Y_S(\mu_D)}{Y_S(\mu_D = 0)}. \quad (7)$$

3. Results and Conclusions

In fig. (2) I show the suppression of the S-wave ground state (1S) for different values of $\mu_D$ as a function of the time duration of the box for both the real and complex potentials. For very small brick lifetimes, less than one period, the suppression is relatively modest for all values of $\mu_D$. For the real potential, the late time suppression saturates to a constant value as long as the in medium potential supports a bounded ground state. When $\mu_D$ is so large that the ground state disappears the late time yield tends to zero. The complex potential, eq. (5), is much more effective in suppressing quarkonia. In this case the late time suppression does not saturate to a constant, but continuously increases due to the finite width of the in-medium states. In fact, at
these late times, the dynamics of the in-medium ground state is dominated by this absorption. It can be checked that after the formation time of the meson the absorption is indeed given by the in-medium width up to an overall shift, which is due to the early time suppression.

In Fig (3) I show the double ratio of the suppression factor of the 2S state to the 1S estate for both potentials. For the real potential, as the temperature rises a modulation in the double yield ratio appears which is due to the non-zero overlap of the vacuum 2S state with the in-medium ground state. This double ratio can be larger than 1, showing a relative enhancement of 2S to 1S states. This is particularly dramatic in the highest temperature case, for which the bound states are loosely bound. This behavior continues at even higher temperatures, when all states are dissolved and it is due to the fact that the vacuum 2S state is wider than the ground state and has a larger overlap with the in-medium continuum. For the complex potential, the late time ratio is constant, which means that both the 2S and 1S yields fall at the same rate. Similarly to the real potential case, this is due to the fact that when the in-medium 2S level is absent, the production is dominated by the overlap with the in-medium 1S. As in the real potential case, the asymptotic value grows with $\mu_D$ and becomes greater than one at larger temperatures.

In summary, by comparing the suppression patterns of a medium-brick characterized by a real quarkonia potential to that of a potential with an imaginary part with the same screening parameter $\mu_D$ I have shown that the non-unitarity effects introduced by the complex potential are very efficient in suppressing quarkonia. Additionally, I have shown that even though the ground state suppression pattern is easy to understand in terms of in-medium bound states, the suppression of excited states, such as 2S-wave level, is due to the mixing with the in-medium 1S state. Such mixing leads to an enhancement of the double ratio of suppression patterns of 2S to 1S states which is due to the wider wave function of the in-medium states. Since this behavior is generic, this type of effects may be behind the anomalous suppression pattern of $\psi(2S)$ states reported by the CMS collaboration [4].

Acknowledgments: I am supported by a Ramón y Cajal fellowship. I also acknowledge financial support by the research grants FPA2010-20807, 2009SGR502 and by the Consolider CPAN project.

[1] J. Casalderrey-Solana, arXiv:1208.2602 [hep-ph].
[2] M. Laine, O. Philipsen, P. Romatschke and M. Tassler, JHEP 0703, 054 (2007) [hep-ph/0611300].
[3] A. Beraudo, J. -P. Blaizot and C. Ratti, Nucl. Phys. A 806, 312 (2008) [arXiv:0712.4394 [nucl-th]].
[4] CMS Collaboration Collaboration, S. Chatrchyan et al.,CMS-PAS-HIN-12-007