Decay width of quarkonia in an anisotropic medium

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Introduction

The experiments at RHIC and LHC have confirmed the formation of quark-gluon plasma (QGP) through the jet quenching, the real or virtual photon spectra, the p_T distribution of secondary hadrons, quarkonia etc. Heavy quarkonium systems have turned out to provide extremely useful probes for the deconfined matter because the force between a heavy quark and its anti-quark, is weakened due to the presence of light quarks and gluons and leads to the dissociation of quarkonium bound states [1]. Based on potential models there were early predictions that J/ψ production would be suppressed in heavy ion collisions. Among the recent theoretical developments in the quarkonium studies, the first-principle calculations of imaginary contributions to the heavy quark potential due to gluonic Landau damping [2], the additional contribution due to singlet to octet transitions etc. [3] are well known. The imaginary part of the potential are generically related to the quarkonium decay processes in the plasma whose consequences on spectral functions [4], thermal widths [2] etc. have recently been studied.
Since the anisotropic distribution is a more realistic description of the parton system generated in heavy-ion collisions, it is worthwhile to consider the properties of quarkonia such as the binding energy, decay width and hence the dissociation in such a system. Traditionally, it was thought that the quarkonium is dissociated when the Debye screening becomes so strong that the corresponding Schrödinger equation does not suppport the bound state solutions. The new suggestion is that the quarkonium effectively dissociates already at a lower temperature [2, 4] when the binding energy is non-zero but overtaken by the Landau-damping induced thermal width [2].

So far we have studied the real part of the potential [5] in the presence of small momentum anisotropy by considering both perturbative as well as the non-perturbative part of the vacuum potential. The imaginary part was calculated earlier by considering only the short-distance part of the vacuum potential, assuming the long-distance part vanishes beyond T_c. But the non-perturbative effect such as the string tension is found to survive till very higher temperatures [6], so we retain the effect of long-distance part in deriving the imaginary part, contrary to others calculation [7] for short distance part. In our work we use the real-time formalism [7] to determine the imaginary part of the potential in the anisotropic medium. So we first obatin the gluon self-energy by folding the approximated phase-space distribution in anisotropic medium ($\xi \ll 1$) as: $f_{\rm aniso}(\mathbf{k}) = f_{\rm iso}\left(\sqrt{\mathbf{k}^2+\xi(\mathbf{k}.\mathbf{n})^2}\right)$ and hence the resummed propagator. The contribution from the quark-loop to the gluon self-energy is of the form [7]

$$\Pi^{\mu\nu}(K) = -\frac{i}{2}N_f g^2 \int \frac{d^4\mathbf{k}}{(2\pi)^4} tr[\gamma^\mu S(Q)\gamma^\nu S(K)] \quad (1)$$

Medium-modification at finite temperature can be obtained by correcting both the short- and long-distance part of the potential (T=0) with a dielectric function $\epsilon(k)$ encoding the effect of deconfinement [8]

$$V(r,T) = \int \frac{d^3\mathbf{k} e^{i\mathbf{k}\cdot\mathbf{r}}}{(2\pi)^{3/2}} \frac{\left(-\sqrt{(\frac{2}{\pi}}\frac{\alpha}{k^2} - \frac{4\sigma}{\sqrt{2\pi}k^4}\right)}{\epsilon(k)} \quad (2)$$

The relation between $\epsilon(k)$ and the 00-component of the gluon propagator is given by

$$\epsilon^{-1}(k) = -\lim_{\omega\to 0} k^2 \Delta^{00}(\omega,k,\xi) \quad (3)$$

In real-time formalism we can calculate the gluon propagator by using the hard-thermal-loop approximation where the longitudinal part of 11-component of the gluon propagator($\Delta_{00}(\omega,k)$) is

$$\Delta^{00}_{11}(\omega,k,\xi) = \frac{1}{2}\left(\Delta^{00}_R + \Delta^{00}_A + \Delta^{00}_F\right), \quad (4)$$

which finally gives the desired potential in the static limit. The real part of the potential, corresponding

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to the short-distance part, thus gives (with $\hat{r} = rm_D$)

$$ReV_1(\mathbf{r},\xi,T) = -\alpha m_D\left[\frac{e^{-\hat{r}}}{\hat{r}} + 1 + \xi\left[\frac{(e^{-\hat{r}}-1)}{6} + e^{-\hat{r}}\left(\frac{1}{6}+\frac{1}{2\hat{r}}+\frac{1}{\hat{r}^2}\right) + \frac{(e^{-\hat{r}}-1)}{\hat{r}^3}\right](1-3\cos^2\theta_r)\right] \quad (5)$$

and the real part of the potential, corresponding to the long-distance part, is

$$ReV_2(\mathbf{r},\xi,T) = \frac{2\sigma}{m_D}\left[\frac{(e^{-\hat{r}}-1)}{\hat{r}} + 1 + 2\xi\left[\frac{(e^{-\hat{r}}+1)}{12} + \frac{(e^{-\hat{r}}-1)}{2\hat{r}}(\frac{2}{3}-\cos^2\theta_r) + \left(\frac{(2e^{-\hat{r}}-1)}{\hat{r}^2} - \frac{(2e^{-\hat{r}}+1)}{12}\right)(1-3\cos^2\theta_r)\right]\right] \quad (6)$$

Now we evalute the imaginary part of the potential where the isotropic part corresponding to the short-distance is

$$ImV_{1(0)}(\mathbf{r},T) = -\alpha T\left[-\frac{\hat{r}^2}{9}(-4+3\gamma_E + 3\log\hat{r})\right] = -\alpha T\phi(\hat{r}) \quad (7)$$

and the anisotropic part is

$$ImV_{1(1)}(\mathbf{r},\xi,T) = (\alpha T\xi)\left[\psi_1^{(1)}(\hat{r},\theta_r) + \psi_1^{(2)}(\hat{r},\theta_r)\right], \quad (8)$$

where the functions $\psi_1^{(1)}(\hat{r},\theta_r)$ and $\psi_1^{(2)}(\hat{r},\theta_r)$ are

$$\psi_1^{(1)}(\hat{r},\theta_r) = \frac{\hat{r}^2}{600}[123 - 90\gamma_E - 90\log\hat{r} + \cos(2\theta_r)(-31+30\gamma_E+30\log\hat{r})]$$
$$\psi_1^{(2)}(\hat{r},\theta_r) = \frac{\hat{r}^2}{90}(-4+3\cos(2\theta_r)) \quad (9)$$

Similarly the isotropic part of string term (with $z = k/m_D$) is.

$$ImV_{2(0)}(r,T) = \frac{2\sigma T}{m_D^2}\phi(\hat{r}) \quad (10)$$

where $\phi(\hat{r})$ is

$$\phi(\hat{r}) = \frac{\hat{r}^2}{6} + \frac{(-107+60\gamma_E+60\log(\hat{r}))\hat{r}^4}{3600} + O(\hat{r}^5) \quad (11)$$

and the anisotropic part is

$$ImV_{2(1)}(r,\xi,T) = -\frac{2\sigma T}{m_D^2}\xi\left[\psi_2^{(1)}(\hat{r},\theta) + \psi_2^{(2)}(\hat{r},\theta)\right] \quad (12)$$

where the functions $\psi_2^{(1)}(\hat{r},\theta)$ and $\psi_2^{(2)}(\hat{r},\theta)$ are given by

$$\psi_2^{(1)}(\hat{r},\theta) = \frac{r^2}{10} + \frac{(-739+420\gamma_E+420\log(\hat{r}))\hat{r}^4}{39200} + \left(-\frac{r^2}{20} + \frac{(176-105\gamma_E-105\log(\hat{r}))\hat{r}^4}{14700})\right)\cos^2\theta_r \quad (13)$$

and

$$\psi_2^{(2)}(\hat{r},\theta) == -\frac{4}{3}\left[\frac{7\hat{r}^2}{120} - \frac{11\hat{r}^4}{3360} + O(\hat{r}^5)\right] - 4\left[-\frac{\hat{r}^2}{60} + \frac{\hat{r}^4}{840} + O(\hat{r}^5)\right]\cos^2\theta_r \quad (14)$$

respectively. The imaginary part of the potential is found to be perturbation to the vacuum potential and thus provides an estimate for the decay width for a particular resonance state. The effects of anisotropy on the imaginary part of the potential and consequently on the thermal width is being investigated further in our work.

References

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