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Imaginary potential and entropic force in non-commutative plasma

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Abstract

We study the imaginary potential and entropic force with respect to a heavy quarkonium in non-commutative $\mathcal{N}=4$ super Yang-Mills (SYM) plasma at strong coupling. We compute the two quantities both along commutative as well as the non commutative coordinates of the brane. It is found that the two methods give the same result: non-commutativity reduces quarkonia dissociation.
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1. Introduction

It is believed that the heavy ion collisions at RHIC and LHC have produced a new state of matter so-called quark gluon plasma (QGP) [1–3]. One of the main experimental signature of QGP formation is the dissociation of quarkonia, such as J/ψ and excited states, in the medium. It was previously suggested that the main mechanism responsible for this suppression is color screening [4]. But recently some authors argued that imaginary potential [5–10] and entropic force [11] can also yield this suppression. However, lots of experiments indicate that QGP is

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strongly coupled, and behaves as a nearly perfect relativistic fluid [12,13]. Therefore, calculational tools for strongly coupled, real time QCD dynamics are needed, such as the Anti-de-Sitter space/conformal field theory (AdS/CFT) duality.

AdS/CFT is a conjectured correspondence between a string theory in AdS space and a conformal field theory in space-time [14–16]. The most well studied example is the duality between the type IIB string theory formulated on $AdS_5 \times S^5$ and $\mathcal{N}=4$ SYM theory in four dimensions. Although $\mathcal{N}=4$ SYM theory differs from QCD in many properties, it reveals some qualitative features of QCD in strongly coupled regime. Already, the AdS/CFT correspondence has yielded many important insights for studying different aspects of QGP [17]. One of the renowned works is the universal value of the ratio between the shear viscosity and the entropy density, $\eta/s = \frac{1}{4\pi}$, for quantum field theories admitting a holographic description [18,19]. Recently, this method has been used to study imaginary potential and entropic force.

The imaginary potential, $\mathrm{ImV}_{Q\bar{Q}}$, could be used to estimate the thermal width which is related to the quarkonia decay process in QGP [20]. Applying AdS/CFT, Noronha and Dumitru have carried out the $\mathrm{ImV}_{Q\bar{Q}}$ for $\mathcal{N}=4$ SYM theory in their seminal work [21]. Therein, the $\mathrm{ImV}_{Q\bar{Q}}$ is related to the effect of thermal fluctuations due to the interactions between the heavy quarks and the medium. Subsequently, this idea has been extended to various cases. For example, the $\mathrm{ImV}_{Q\bar{Q}}$ of static quarkonia is studied in [22]. The $\mathrm{Im}V_{Q\bar{Q}}$ of moving quarkonia is discussed in [23,24]. Also, this quantity has been investigated in strongly coupled anisotropic plasma [25] as well as some AdS/QCD models [26,27]. Besides, there are other ways to study $\mathrm{Im}V_{Q\bar{Q}}$ from holography [28–30].

Another quantity sensitive to the quakonia dissociation is the entropic force. This force was introduced in [31] to explain the elasticity of polymer strands in rubber. Later, Verlinde proposed [32] that it would be responsible for gravity (but this idea may be controversial [33] and will not be discussed here). Recently, D. E. Kharzeev [11] argued that this force can drive the dissociation process if one considers the process of deconfinement as an entropic self-destruction. His argument is based on the Lattice results which indicate that there is a peak in the heavy quark entropy around the crossover region of the QGP [34–36]. In a more recent work [37], K. Hashimoto et al. were among the first to carry out the entropic force of static quarkonia from AdS/CFT. It is found that the peak of the entropy near the transition point associates with the nature of deconfinement. Subsequently, K. B. Fadafan et al. have studied the entropic force of moving quarkonia in [38]. Further study in this direction can be found in [39–42].

Here we extend the study of the imaginary potential and entropic force in non-commutative Yang-Mills (NCYM) plasma. It is known that the gauge theory develops a space-space non-commutativity when Dp-brane world-volume is subjected to a large asymptotic magnetic or B-field [43–45]. Also, understanding holography in NCYM theory is by itself an interesting endeavor. Already, there are several quantities or observables that have been studied in such theories. For example, the shear viscosity to entropy density ratio has been discussed in [46] and it is shown show that the longitudinal η/s differs significantly from that of its value from the transverse fluctuations, in particular, one could have a natural violation of the conjectured lower bound on η/s solely from the anisotropic considerations even in Einstein gravity. Also, the drag force was studied in [47] and it is found that when the quark moves in one of the NC directions the drag force gets reduced and thus non-commutativity makes the plasma less viscous. Moreover, the hydrodynamics of charge diffusion was analyzed in [48] and it turns out that the DC conductivity gets significantly modified along the NC directions of the brane and its value is lower than that of the commutative counterpart. Motivated by this, in this paper we study the imaginary potential and entropic force in NCYM plasma. Specifically, we want to see how non-commutativity affects the two quantities as well as the quarkonia dissociation. This is the purpose of the present work.

The organization of the paper is as follows. In the next section, we introduce the geometrical construction in the dual gravitational counterpart of the non-commutative given in [46]. In section 3, we study the imaginary potential of quarkonia in this background and discuss how non-commutativity affects it. In section 4, we analyze the entropic force of heavy quarkonia for the same background. The last part is devoted to conclusion and discussion.

2. Setup

Let us begin with a briefly review of the geometrical construction in the bulk space time which is holographically dual to non-commutative $\mathcal{N}=4$ SYM theory at strong coupling. It is known that [43–45] non-commutative gauge theories at strong coupling can be obtained from string theory by considering the decoupling limit in a system of D_p branes in the presence of a background NS B field, which gives rise to certain scale of non-commutativity in the large N limit. The non commutative $\mathcal{N}=4$ SYM theory at finite temperature whose dual counterpart in the string frame is given by [46]

$$ds^2 = H^{-1/2}(-f dt^2 + dx^2 + h(dy^2 + dz^2)) + H^{1/2}(\frac{dr^2}{f} + r^2 d\Omega_5^2), \tag{1}$$

with

$$f = 1 - \frac{r_h^4}{r^4}, \qquad h = \frac{1}{1+\Theta^2 H^{-1}}, \qquad H = \frac{R^4}{r^4}, \tag{2}$$

where R is the AdS radius. r describes the 5th dimensional coordinate with $r = r_h$ the horizon and $r = \infty$ the boundary. Θ represents the non-commutative parameter. So (t, x) are the usual commutative directions while (y, z) exhibit the non-commutative nature. Also, it should be noticed that due to the presence of the non commutativity along (y, z) of the brane, the full SO(3) symmetry of the boundary theory reduces to SO(2), leaving the rotational invariance only over the (y, z) plane.

In addition, the Hawking temperature reads

$$T = \frac{r_h}{\pi R^2}, \tag{3}$$

one can see that T is independent of Θ.

3. Imaginary potential

In this section we follow the argument in [21] to study the imaginary potential for the background metric (1). First we consider the quark anti-quark pair located at one of the non-commutative directions, e.g., y direction. In this case, the coordinates are parameterized as

$$t = \tau, \qquad x = 0, \qquad y = \sigma, \qquad z = 0, \qquad r = r(\sigma). \tag{4}$$

The Nambu-Goto action for the string world sheet is

$$S = -\frac{1}{2\pi\alpha'}\int d\tau d\sigma \mathcal{L} = -\frac{1}{2\pi\alpha'}\int d\tau d\sigma \sqrt{-g}, \tag{5}$$

with g the determinant of the induced metric and

$$g_{\alpha\beta} = g_{\mu\nu}\frac{\partial X^\mu}{\partial\sigma^\alpha}\frac{\partial X^\nu}{\partial\sigma^\beta}, \tag{6}$$

where $g_{\mu\nu}$ and X^μ are the metric and target space coordinates, respectively. According to the AdS/CFT dictionary, $\lambda \equiv g^2_{\rm YM}N = \frac{R^4}{\alpha'^2}$ with N the number of D3 branes and λ the 't Hooft coupling.

Substituting (4) into (1), the Lagrangian density in (5) reads

$$\mathcal{L} = \sqrt{\frac{r^4 fh}{R^4} + (\frac{dr}{d\sigma})^2}. \tag{7}$$

Since the Lagrangian does not depend on σ explicitly, the Hamiltonian satisfies

$$\mathcal{L} - \frac{\partial\mathcal{L}}{\partial\frac{dr}{d\sigma}}\frac{dr}{d\sigma} = constant. \tag{8}$$

Imposing the boundary condition of the deepest point

$$\frac{dr}{d\sigma} = 0, \qquad r = r_c, \tag{9}$$

one gets

$$\frac{dr}{d\sigma} = \sqrt{\frac{a^2(r) - a(r)a(r_c)}{a(r_c)}}, \tag{10}$$

with

$$a(r) = \frac{r^4 fh}{R^4},$$
$$a(r_c) = \frac{r_c^4 f(r_c)h(r_c)}{R^4},$$
$$f(r_c) = 1 - (\frac{r_h}{r_c})^4,$$
$$h(r_c) = \frac{1}{1+\Theta^2 r_c^4/R^4}. \tag{11}$$

For later convenience, we set $R = 1$. Integrating (10), the inter-distance of $Q\bar{Q}$ reads

$$L = 2\int_{r_c}^{\infty} dr\sqrt{\frac{a(r_c)}{a^2(r) - a(r)a(r_c)}} = 2\int_{r_c}^{\infty} dr\sqrt{\frac{1}{\frac{r^4-r_h^4}{1+\Theta^2 r^4}(\frac{r^4-r_h^4}{r_c^4-r_h^4}\frac{1+\Theta^2 r_c^4}{1+\Theta^2 r^4} - 1)}}, \tag{12}$$

from the above equation one finds that increasing Θ leads to increasing L. Namely, the non-commutative effects increase the value of inter-distance from that of its usual value corresponding to the commutative case.

Next, we calculate the heavy quark potential. The real part of the potential is given by [49–51]

$$ReV_{Q\bar{Q}} = \frac{1}{\pi\alpha'}\int_{r_c}^{\infty} dr[\sqrt{\frac{a(r)}{a(r) - a(r_c)}} - 1] - \frac{1}{\pi\alpha'}\int_{r_h}^{r_c} dr. \tag{13}$$

On the other hand, the imaginary part of the potential can be obtained by the thermal world-sheet fluctuation method [21], given by

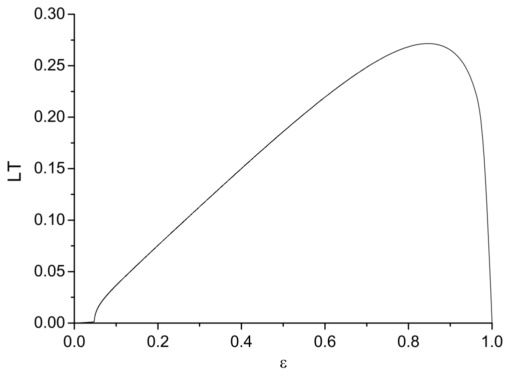

Fig. 1. LT versus ε with $\Theta = 0.001$ GeV^{-2}.

$$ImV_{Q\bar{Q}} = -\frac{1}{2\sqrt{2}\alpha'}[\frac{a'(r_c)}{2a''(r_c)} - \frac{a(r_c)}{a'(r_c)}], \tag{14}$$

where

$$\begin{aligned}
a'(r_c) &= r_c^4 f(r_c)h'(r_c) + r_c^4 f'(r_c)h(r_c) + 4r_c^3 f(r_c)h(r_c), \\
a''(r_c) &= r_c^4 f(r_c)h''(r_c) + r_c^4 f''(r_c)h(r_c) + 2r_c^4 f'(r_c)h'(r_c) + 8r_c^3 f(r_c)h'(r_c) \\
&\quad + 8r_c^3 f'(r_c)h(r_c) + 12r_c^2 f(r_c)h(r_c), \\
h'(r_c) &= -\frac{4\Theta^2 r_c^3}{(1+\Theta^2 r_c^4)^2}, \qquad h''(r_c) = \frac{32\Theta^4 r_c^6 - 12\Theta^2 r_c^2(1+\Theta^2 r_c^4)}{(1+\Theta^2 r_c^4)^3}, \\
f'(r_c) &= 4r_h^4 r_c^{-5}, \qquad f''(r_c) = -20 r_h^4 r_c^{-6},
\end{aligned} \tag{15}$$

note that when Θ is put to zero, the result of ordinary $\mathcal{N} = 4$ SYM plasma [21] is recovered.

Before going further, we discuss the value range of Θ. There are various disparate experimental bounds on Θ from different physical considerations, e.g., the range is claimed to be $\sim (1 - 10 \text{ TeV})^{-2}$ [52], $\sim (10^{12} - 10^{13} \text{ GeV})^{-2}$ [53], or $\sim (10^{15} \text{ GeV})^{-2}$ [54]. One can see that in all these cases the values of Θ are very small, so there is little hope of getting a significant correction due to non-commutativity in current experiments [46,47]. However, it was argued [43–45] that the non-commutativity introduces a non-locality in space due to space uncertainly and these non-local effects will enhance as the temperature increases. For that reason, the NC effect may be observed at forthcoming high energy collider experiments. Here for the purpose of calculating convenience, we choose $\Theta = 0.001, 0.005$ GeV^{-2} in the numerics.

In addition, there are several restrictions on this model. First, the imaginary potential should be negative, yielding

$$\frac{a'(r_c)}{2a''(r_c)} - \frac{a(r_c)}{a'(r_c)} > 0, \tag{16}$$

results in

$$\varepsilon > \varepsilon_{min}, \tag{17}$$

with $\varepsilon \equiv r_h/r_c$, where ε_{min} can be obtained numerically.

The second restriction is in relation to the maximum value of L. To illustrate this, we plot LT as a function of ε for $\Theta = 0.001$ GeV^{-2} in Fig. 1. Other cases with different values of Θ have

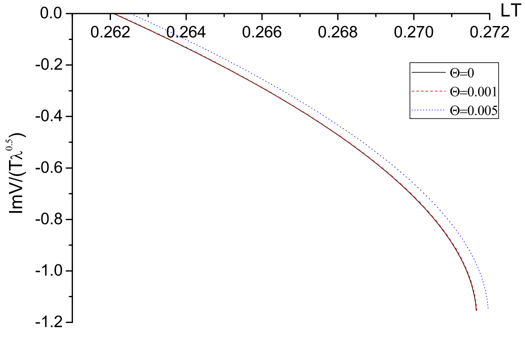

Fig. 2. Im$V/(\sqrt{\lambda}T)$ against LT. From left to right $\Theta = 0, 0.001, 0.005$ GeV^{-2}, respectively.

similar picture. One can see that LT increases firstly and then decreases with the increase of ε. In fact, for the later case of $\varepsilon > \varepsilon_{max}$ (ε_{max} denotes the value of ε at $LT = LT_{max}$), one needs to consider highly curved configurations for the string worldsheet which are not solutions of the Nambu-Goto action [55]. Here we are more interested in the former, corresponding to $\varepsilon < \varepsilon_{max}$. Taken together, the domain of applicability of the model is $\varepsilon_{min} < \varepsilon < \varepsilon_{max}$.

In Fig. 2, we plot Im$V/(\sqrt{\lambda}T)$ versus LT for $\Theta = 0, 0.001, 0.005$ GeV^{-2}, respectively. One can see that for each plot the Im$V_{Q\bar{Q}}$ starts at a L_{min}, corresponding to $\varepsilon = \varepsilon_{min}$, and ends at a L_{max}, corresponding to $\varepsilon = \varepsilon_{max}$. In addition, increasing Θ, Im$V_{Q\bar{Q}}$ is generated for larger distance. As discussed in [23], the dissociation properties of quarkonia should be sensitive to Im$V_{Q\bar{Q}}$, and if the onset of Im$V_{Q\bar{Q}}$ happens for smaller LT, the quarkonia suppression will be stronger. Since Θ increases the onset of LT, one concludes that the inclusion of non-commutativity reduces quarkonia dissociation. Interestingly, similar observations have also been made earlier in a different context of anisotropy [25].

Analogously, one can analyze the quark anti-quark pair located at the commutative direction, i.e., x direction, it is found that the imaginary potential in this case does not get modified due to the presence of the NC parameter.

4. Entropic force

In this section, we follow the calculations of [37] to study the entropic force for the background (1). The entropic force is defined as [11]

$$\mathcal{F} = T\frac{\partial \mathcal{S}}{\partial L}, \tag{18}$$

where $\mathcal{S}$ denotes the entropy. L and T represent the inter-quark distance and temperature, respectively.

As before, one first considers the quark antiquark pair located at one of the non-commutative directions, e.g., y direction. Parallel to the previous section, one gets

$$L = 2\int_{r_c}^{\infty} dr\sqrt{\frac{A(r_c)}{A^2(r) - A(r)A(r_c)}}, \tag{19}$$

with

$$A(r) = r^4 f h,$$
$$A(r_c) = r_c^4 f(r_c) h(r_c),$$
$$f(r_c) = 1 - (\frac{r_h}{r_c})^4,$$
$$h(r) = \frac{1}{1 + \Theta^2 r^4},$$
$$h(r_c) = \frac{1}{1 + \Theta^2 r_c^4}. \tag{20}$$

What follows is to calculate the entropy, given by

$$\mathcal{S} = -\frac{\partial F}{\partial T}, \tag{21}$$

where F is the free energy of $Q\bar{Q}$. There are two options for F.

1. If $L > \frac{c}{T}$ (where c denotes the maximum value of LT and it could be determined numerically), one needs to take into account new configurations [55]. Given that, the choice of the free energy $F^{(1)}$ is not unique [56]. Here we choose a configuration of two disconnected trailing drag strings [57,58],

$$F^{(1)} = \frac{1}{\pi\alpha'} \int_{r_h}^{\infty} dr. \tag{22}$$

By virtue of (21), one gets

$$\mathcal{S}^{(1)} = \sqrt{\lambda}\theta(L - \frac{c}{T}), \tag{23}$$

where $\theta(L - \frac{c}{T})$ represents the Heaviside step function.

2. If $L < \frac{c}{T}$, the free energy could be derived from the on-shell action of the fundamental string in the dual geometry,

$$F^{(2)} = \frac{1}{\pi\alpha'} \int_{r_c}^{\infty} dr \sqrt{\frac{A(r)}{A(r) - A(r_c)}}. \tag{24}$$

Using (21) and (24), one obtains

$$\mathcal{S}^{(2)} = -\frac{1}{2\alpha'} \int_{r_c}^{\infty} dr \frac{A(r)A'(r_c) - A'(r)A(r_c)}{\sqrt{A(r)}[A(r) - A(r_c)]^{3/2}}, \tag{25}$$

with

$$A'(r) = -4 r_h^3 h,$$
$$A'(r_c) = -4 r_h^3 h(r_c), \tag{26}$$

note that for $\Theta = 0$ in (25) the result of [37] is reproduced.

To study the effect of the non-commutativity on the entropic force, we plot $S^{(2)}/\sqrt{\lambda}$ versus LT for $\Theta = 0, 0.001, 0.005$ GeV^{-2} in Fig. 3, respectively (where we have used the relation $\sqrt{\lambda} = R^2/\alpha' = 1/\alpha'$). Note that the right panel is a part of the left one. From the right one, one can see that increasing Θ leads to smaller entropy at small distances. As mentioned earlier, the growth

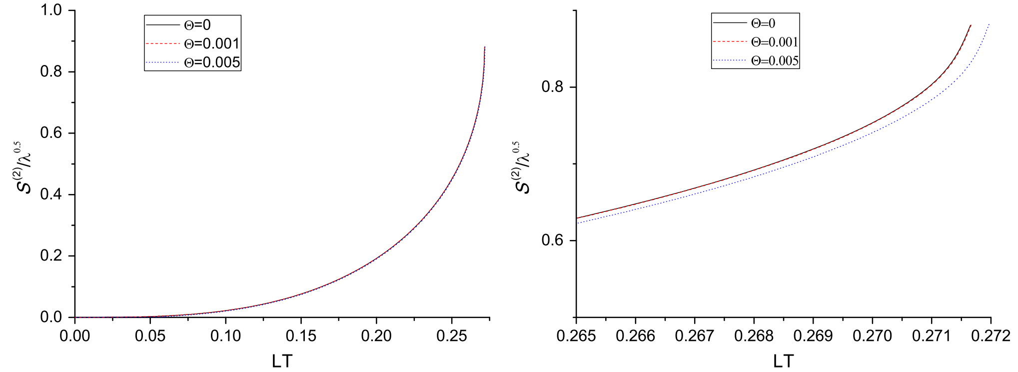

Fig. 3. $\mathcal{S}^{(2)}/\sqrt{\lambda}$ versus LT. Left: complete graph. Right: local image. In the plots from top to bottom $\Theta = 0, 0.001, 0.005$ GeV^{-2}, respectively.

of the entropy with the inter-distance is responsible for the entropic force which destructs the quarkonia. Therefore, one concludes that the presence of non-commutativity tends to decrease the entropic force thus reducing quarkonia dissociation, in agreement with the analysis of the imaginary potential.

Likewise, if one considers the quark anti-quark pair at the commutative direction, one will find that the entropic force is not modified by the NC effect.

5. Conclusion

Imaginary potential and entropic force may represent two different mechanisms for melting the heavy quarkonia. In this paper, we investigated the two quantities in NCYM plasma at strong coupling. The motivation rests on the earlier studies on η/s [46] and drag force [47] in such theories. It is shown that the imaginary potential and entropic force get modified (only) along the NC directions of the brane. In particular, increasing Θ, the imaginary potential is generated for larger distance thus decreasing quarkonia dissociation. While increasing Θ leads to decreasing the entropic force thus reducing quarkonia dissociation. Namely, the two methods achieve the same results: non-commutativity reduces quarkonia dissociation.

On the other hand, the anisotropy in the QGP can also be related to an external magnetic field, given its creation during the early stages of noncentral heavy ion collisions [59–63], and recently the imaginary potential has been studied in $\mathcal{N} = 4$ SYM plasma under the influence of magnetic field [64]. It was shown that the presence of magnetic field enhances quarkonia dissociation, reverse to the effect of non-commutativity. Similar findings appear in the study of drag force, i.e., the NC effects decrease the drag force [47] but the magnetic field enhances it [65–67]. Those disagreements are conceivable, because the sources of anisotropy of those models are different.

Finally, it would be interesting to study the imaginary potential and entropic force in other anisotropic version of $\mathcal{N} = 4$ SYM theories, e.g., θ deformed theory [68] and compare the results of those models with this work as well as [64]. We leave this as a future study.

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