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Back reaction effects on the imaginary potential of quarkonia in heavy quark cloud

Ping-ping Wu [a], Xiangrong Zhu [b], Zi-qiang Zhang [c,*]

[a] *School of Physics and Electronic Information Engineering, Henan Polytechnic University, Jiaozuo 454000, China*
[b] *School of Science, Huzhou University, Huzhou 313000, China*
[c] *School of Mathematics and Physics, China University of Geosciences, Wuhan 430074, China*

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Abstract

Applying the AdS/CFT correspondence, we investigate the effect of back reaction on the imaginary part of heavy quarkonia potential in strongly coupled $\mathcal{N} = 4$ supersymmetric Yang-Mills (SYM) plasma. The back reaction considered here arises from the inclusion of static heavy quarks uniformly distributed over $\mathcal{N} = 4$ SYM plasma. It is shown that the presence of back reaction reduces the absolute value of the imaginary potential thus decreasing the thermal width. Furthermore, the results imply that back reaction enhances the quarkonia dissociation.

1. Introduction

The heavy ion collisions at RHIC and LHC are believed to produce a new phase of matter so-called quark gluon plasma (QGP) [1]. One of the main experimental signatures of QGP formation is quarkonia dissociation. Early evidence [2] indicated that the binding interaction of the quark-antiquark pair ($Q\bar{Q}$) is screened by the medium, causing this dissociation. However, recent research [3–8] suggests that the imaginary part of a quarkonia potential, $\mathrm{ImV}_{Q\bar{Q}}$, may be a

* Corresponding author.
E-mail addresses: wupp@hpu.edu.cn (P.-p. Wu), xrongzhu@zjhu.edu.cn (X. Zhu), zhangzq@cug.edu.cn (Z.-q. Zhang).

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more important reason than screening. Also, this quantity could be used to evaluate the thermal width of quarkonia. Already, the ImV_{QQ} has been studied in perturbative QCD etc. [9–11]. However, lots of experiments show that QGP is strongly coupled [12], which involves nonperturbative physics suitable for the Anti-de-Sitter space/conformal field theory (AdS/CFT) correspondence.

AdS/CFT [13–15], namely the duality between the type IIB string theory on $AdS_5 \times S^5$ and $\mathcal{N}=4$ SYM with an SU(N) gauge group in four dimensions, provides a powerful tool to deal with strongly coupled field systems. During the last two decades, this duality has offered many important insights for studying various aspects of QGP (see [16] for a good review). Using AdS/CFT, Noronha and Dumitru have calculated the ImV_{QQ} for $\mathcal{N}=4$ SYM plasma in their seminal work in [17], where they argued that the ImV_{QQ} could be obtained by analyzing the effect of thermal fluctuations due to the interactions between the heavy quarks and the medium. Subsequently, this idea has been generalized to various cases. For instance, the Im$V_{Q\bar{Q}}$ of static quarkonia is studied in [18,19]. The Im$V_{Q\bar{Q}}$ of moving quarkonia is addressed in [20,21]. The finite 't Hooft coupling corrections on Im$V_{Q\bar{Q}}$ is achieved in [22]. The chemical potential effect on Im$V_{Q\bar{Q}}$ is discussed in [23]. For a study of Im$V_{Q\bar{Q}}$ from AdS/QCD, see [24–26]. Moreover, this quantity has been holographically computed in some other ways [27–29].

The purpose of this paper is to study the effect of back reaction on Im$V_{Q\bar{Q}}$ for $\mathcal{N}=4$ SYM plasma. As we know, QGP is comprised of huge amounts of free quarks and gluons. That would mean if one analyzes the dynamics of a heavy quark (or quarkonium) in this hot plasma, the back reaction of the plasma should be taken into account. To our knowledge, except for a few examples [30,31] etc., it remains very difficult to investigate the strongly coupled gauge theory with large number of flavour quarks. Recently, S. Chakrabortty argued [32] that the back reaction could be modeled by means of the deformation of the geometry due to finite density string cloud. To be specific, a back reacted gravity background could be dual to a system of large number of heavy quarks uniformly distributed over the $\mathcal{N}=4$ SYM plasma. Later, the drag force [32] and jet quenching parameter [33] have been studied in such theories. It was shown that the presence of back reaction increases the two quantities thus enhancing the energy loss. Motivated by this, in this paper we would like to study the imaginary potential in such theories. We want to know: how back reaction modifies the imaginary potential and the thermal width? One step further, how back reaction affects the quarkonia dissociation?

The structure of the paper is as follows. In the next section, we introduce the backreacted gravity background given in [32]. In section 3, we evaluate the expressions for the real and imaginary part of the quarkonia potential, in turn. In section 4, we analyze the effect of back reaction on the imaginary potential and discuss how such effects modify the thermal width as well as the quarkonia dissociation. In the last section, we summarize our results and make some discussions.

2. Background geometry

One considers the $(n+1)$-dimensional gravitational action [32]

$$S=\frac{1}{16\pi G_{n+1}}\int d^5x\sqrt{-g}(\mathcal{R}-2\Lambda)+S_m, \tag{1}$$

where G_{n+1} is the $(n+1)$-dimensional Newton constant. g denotes the determinant of the metric g_{MN}. $\mathcal{R}$ represents the Ricci scalar. Λ refers to the cosmological constant. S_m stands for the matter part, given by

$$S_m = -\frac{1}{2}\sum_i \mathcal{T}_i \int d^2\xi \sqrt{-h} h^{\alpha\beta} \partial_\alpha X^\mu \partial_\beta X^\nu g_{\mu\nu}, \tag{2}$$

where $\mathcal{T}_i$ is the tension. $h^{\alpha\beta}$ and $g^{\mu\nu}$ are the world-sheet metric and space time metric with α, β the world sheet coordinates and μ, ν the space time directions.

Varying action (1) with respect to the space time metric yields

$$\mathcal{R}_{\mu\nu} - \frac{1}{2}\mathcal{R} g_{\mu\nu} + \Lambda g_{\mu\nu} = 8\pi G_{n+1} T_{\mu\nu}, \tag{3}$$

with

$$T^{\mu\nu} = -\sum_i \mathcal{T}_i \int d^2\xi \frac{1}{\sqrt{|g_{\mu\nu}|}} \sqrt{|h_{\alpha\beta}|} h^{\alpha\beta} \partial_\alpha X^\mu \partial_\beta X^\nu \delta_i^{n+1}(x - X), \tag{4}$$

where the delta function denotes the source divergences due to the presence of the strings.

To proceed, one considers the space-time metric of the form

$$ds^2 = g_{tt}(r)dt^2 + g_{rr}(r)dr^2 + r^2 \delta_{ab} dx^a dx^b, \tag{5}$$

where (a, b) run over $(n-1)$ space direction.

Using the static gauge $t = \xi^0$, $r = \xi^1$. The non vanishing components of $T^{\mu\nu}$ become

$$T^{tt} = -\frac{b g^{tt}}{r^{n-1}}, \qquad T^{rr} = -\frac{b g^{rr}}{r^{n-1}}, \tag{6}$$

where one have assumed the strings are uniformly distributed over $(n-1)$ direction. Given that, the density reads

$$b(x) = T \sum_i \delta_i^{n-1}(x - X_i), \qquad with \qquad b > 0, \tag{7}$$

note that for $b < 0$, $T^{\mu\nu}$ will no longer satisfy the weak and the dominant energy conditions. One looks for a solution of (3) in AdS space and parametrizes the metric accordingly treating b as a constant

$$ds^2 = -V(r)dt^2 + \frac{dr^2}{V(r)} + r^2 h_{ij} dx^i dx^j, \tag{8}$$

where h_{ij} represents the metric on the $(n-1)$-dimensional boundary. Since for the matter part one will mainly focus on the string cloud for which the non vanishing components of $T^{\mu\nu}$ are

$$T_t^t = T_r^r = -\frac{b}{r^{n-1}}, \qquad with \qquad b > 0. \tag{9}$$

By solving the Einstein's equation, one has

$$V(r) = K + \frac{r^2}{R^2} - \frac{2m}{r^{n-2}} - \frac{2b}{(n-1)r^{n-3}}, \tag{10}$$

where R denotes the AdS radius. $K = 0, -1, 1$ depending on whether the boundary is flat, spherical or hyperbolic respectively. In this work, we are mostly interested in the $K = 0, n = 4$ case and the corresponding metric is

$$ds^2 = \frac{r^2}{R^2}(-f(r)dt^2 + d\vec{x}^2) + \frac{R^2}{r^2 f(r)} dr^2, \tag{11}$$

with

$$f(r)=1-\frac{2mR^2}{r^4}-\frac{2}{3}\frac{bR^2}{r^3}, \tag{12}$$

where r denotes the coordinate describing the 5th dimension with $r=\infty$ the boundary. The event horizon r_h satisfies $f(r_h)=0$. Given that, one can write m as

$$m=(1-\frac{2}{3}\frac{bR^2}{r_h^3})\frac{r_h^4}{2R^2}. \tag{13}$$

Moreover, the temperature is

$$T=\frac{\sqrt{g^{rr}}\partial_r\sqrt{g_{tt}}}{2\pi}\Big|_{r=r_h}=\frac{6r_h^3-bR^2}{6\pi R^2 r_h^2}. \tag{14}$$

Note that the geometry (11) is thermodynamically stable under tensor and vector perturbations. Also, it resembles AdS Schwarzschild black hole with negative curvature horizon. For more details about the backreacted gravity background, we refer to [32,33].

3. ImV$_{Q\bar{Q}}$ for the back reacted gravity background

In this section, we follow the approach in [17,18] to evaluate the expressions for the real and imaginary part of the quarkonia potential for the background metric (11). The expectation value of the static (temporal) Wilson loop is given by [34]

$$W(C)=\frac{1}{N_c}TrPe^{i\int A_\mu dx^\mu}, \tag{15}$$

where C denotes a closed loop in a 4-dimensional space time and the trace is over the fundamental representation of the $SU(N_c)$ group. A_μ represents the gauge potential. P enforces the path ordering along C. Considering a rectangular loop with one direction along the time coordinate $\mathcal{T}$ and spatial extension L. Then in the asymptotic limit $\mathcal{T}\to\infty$, the heavy quark potential $V_{Q\bar{Q}}(L)$ can be extracted from the vacuum expectation of the rectangular Wilson loop,

$$<W(C)>\sim e^{-i\mathcal{T}V_{Q\bar{Q}}(L)}. \tag{16}$$

On the other hand, in the supergravity limit,

$$<W(C)>\sim e^{-iS_{str}}, \tag{17}$$

with S_{str} the classical Nambu-Goto action of a string in the bulk,

$$S_{str}=S=-\frac{1}{2\pi\alpha'}\int d\tau d\sigma\sqrt{-det(G_{\mu\nu}\partial_a X^\mu\partial_b X^\nu)}, \tag{18}$$

where $X^\mu(\sigma,\tau)$ are the worldsheet embedding coordinates with $\mu,\nu=0,1,...,4$ and $a,b=\sigma,\tau$. α' is related to 't Hooft coupling λ as $R^2/\alpha'=\sqrt{\lambda}$.

To proceed, one uses the remaining symmetry of (18) to completely fix the static gauge, such that

$$t=\tau,\qquad x=\sigma,\qquad y=0,\qquad z=0,\qquad r=r(\sigma). \tag{19}$$

Under this assumption, the string action with end points fixed at $x = \pm L/2$ reads

$$S = \frac{\mathcal{T}}{2\pi\alpha'} \int_{-L/2}^{L/2} d\sigma \sqrt{\frac{r^4 f(r)}{R^4} + (\frac{dr}{d\sigma})^2}. \tag{20}$$

Since the above action does not depend on σ explicitly, the corresponding Hamiltonian is a constant

$$\mathcal{L} - \frac{\partial \mathcal{L}}{\partial(\frac{dr}{d\sigma})} \frac{dr}{d\sigma} = \frac{r^4 f(r)/R^4}{\sqrt{\frac{r^4 f(r)}{R^4} + (\frac{dr}{d\sigma})^2}} = constant. \tag{21}$$

Imposing the boundary condition at the deepest point of the U-shaped string like

$$\frac{dr}{d\sigma} = 0, \qquad r = r_*, \tag{22}$$

one gets

$$\frac{dr}{d\sigma} = \sqrt{\frac{a^2(r) - a(r)a(r_*)}{a(r_*)}}, \tag{23}$$

where

$$a(r) = \frac{r^4 f(r)}{R^4}, \qquad a(r_*) = \frac{r_*^4 f(r_*)}{R^4}, \tag{24}$$

with

$$f(r_*) = 1 - \frac{2mR^2}{r_*^4} - \frac{2}{3}\frac{bR^2}{r_*^3}. \tag{25}$$

Integrating (23), the inter-distance of $Q\bar{Q}$ can be written as

$$L = 2\int_{r_*}^{\infty} dr \frac{d\sigma}{dr} = 2\int_{r_*}^{\infty} dr \sqrt{\frac{a(r_*)}{a^2(r) - a(r)a(r_*)}}. \tag{26}$$

Putting (23) into (20), the total action for $Q\bar{Q}$ is obtained as

$$S = \frac{\mathcal{T}}{\pi\alpha'} \int_{r_*}^{\infty} dr \sqrt{\frac{a(r)}{a(r) - a(r_*)}}. \tag{27}$$

The above action is divergent since it contains the self-energies of $Q\bar{Q}$ pair. One could cure this divergence by subtracting from S the self energy contribution of the two quarks [34–36]. Then the real part of the heavy quark potential for the background metric (11) reads

$$ReV_{Q\bar{Q}} = \frac{1}{\pi\alpha'} \int_{r_*}^{\infty} dr [\sqrt{\frac{a(r)}{a(r) - a(r_*)}} - 1] - \frac{1}{\pi\alpha'} \int_{r_h}^{r_*} dr. \tag{28}$$

For the above equation, we would like to make the following comment. For small distance L, one expects T as well as b, to have negligible effect on the interaction potential of $Q\bar{Q}$. This

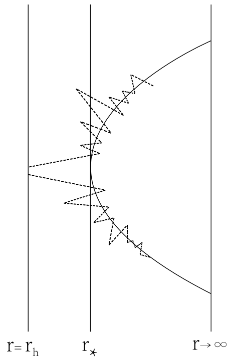

Fig. 1. The effect of thermal fluctuations (dashed line) around the classical configuration (solid line).

argument is backed by lattice results [37,38] which indicates that for $LT << 1$ the free energy is independent of T. Furthermore, for small L if $Q\bar{Q}$ interaction does not depend on T and b, then the regularization in (28) does not depend on these scales either.

To proceed, we derive the expressions of $\mathrm{ImV}_{Q\bar{Q}}$. One considers the effect of thermal world sheet fluctuations around the classical configurations $r_c(x)$,

$$r(x) = r_c(x) \to r(x) = r_c(x) + \delta r(x), \tag{29}$$

where $r_c(x)$ solves $\delta S_{NG} = 0$. For simplicity, $\delta r(x)$ is taken to be of arbitrarily long wavelength, i.e., $\frac{d\delta r(x)}{dx} \to 0$. The physical picture of the thermal fluctuations is shown in Fig. 1. Notice that if r_* is close enough to r_h, the fluctuations of very long wavelength may reach the horizon.

The string partition function that takes into account the fluctuations is given by

$$Z_{str} \sim \int \mathcal{D}\delta r(x) e^{iS_{NG}(r_c(x)+\delta r(x))}. \tag{30}$$

Discretizing the interval $-L/2 < x < L/2$ into $2N$ points located at $x_j = j\Delta x$ ($j = -N, -N+1, ..., N$) with $\Delta x \equiv L/(2N)$, one has

$$Z_{str} \sim \lim_{N\to\infty} \int d[\delta r(x_{-N})]\cdots d[\delta r(x_N)] exp[\frac{i\mathcal{T}\Delta x}{2\pi\alpha'} \sum_j \sqrt{(r_j')^2 + a(r_j)}], \tag{31}$$

where $r_j \equiv r(x_j)$ and $r_j' \equiv r'(x_j)$. The thermal fluctuations are most important around $x = 0$, where $r = r_*$. Hence, one can expand $r_c(x_j)$ around $x = 0$, keeping only terms up to second order in x_j. Given that $r_c'(0) = 0$, one gets

$$r_c(x_j) \approx r_* + \frac{x_j^2}{2} r_c''(0). \tag{32}$$

Also, the expansion for $a(r_j)$, keeping only terms up to second order in $x_j^m \delta r_n$, reads

$$a(r_j) \approx a_* + \delta r a_*' + r_c''(0) a_*' \frac{x_j^2}{2} + \frac{\delta r^2}{2} a_*'', \tag{33}$$

with $a_* \equiv a(r_*)$, $a_*' \equiv a'(r_*)$, etc. Then, the exponent in (31) can be approximated as

$$S_j^{NG} = \frac{\mathcal{T}\Delta x}{2\pi\alpha'} \sqrt{C_1 x_j^2 + C_2}, \tag{34}$$

where

$$C_1 = \frac{r_c''(0)}{2}[2r_c''(0) + a_*'], \qquad C_2 = a_* + \delta r a_*' + \frac{\delta r^2}{2} a_*''. \tag{35}$$

If the function in the square root of (34) is negative, S_j^{NG} will contribute to $\mathrm{ImV}_{QQ} \neq 0$. The relevant region of the fluctuations is the one between δr that yields a vanishing argument in the square root of (34). Thus, one could isolate the $j-$th contribution

$$I_j \equiv \int_{\delta r_{jmin}}^{\delta r_{jmax}} d(\delta r_j) exp[\frac{i\mathcal{T}\Delta x}{2\pi\alpha'}\sqrt{C_1 x_j^2 + C_2}], \tag{36}$$

where δr_{jmin} and δr_{jmax} are the roots of $C_1 x_j^2 + C_2$ in δr.

The integral in (36) can be obtained by using the saddle point method in the classical gravity approximation, i.e., $\alpha' << 1$. The exponent has a stationary point when the function inside the root square of (36)

$$D(\delta r_j) \equiv C_1 x_j^2 + C_2(\delta r_j), \tag{37}$$

assumes an extremal value. This happens for

$$\delta r = -\frac{a_*'}{a_*''}. \tag{38}$$

Requiring that the square root has an imaginary part, results in

$$D(\delta r_j) < 0 \rightarrow -x_c < x_j < x_c, \tag{39}$$

where

$$x_c = \sqrt{\frac{1}{C_1}(\frac{a_*'^2}{2a_*''} - a_*)}. \tag{40}$$

Taking $x_c = 0$ if the square root in (40) is not real. With these conditions, one can approximate $D(\delta r)$ by $D(-\frac{a_*'}{a_*''})$ in (36)

$$I_j \sim exp[\frac{i\mathcal{T}\Delta x}{2\pi\alpha'}\sqrt{C_1 x_j^2 + a_* - \frac{a_*'^2}{2a_*''}}]. \tag{41}$$

Then the total contribution to the imaginary part is given by $\Pi_j I_j$, leading to

$$ImV_{Q\bar{Q}} = -\frac{1}{2\pi\alpha'}\int_{|x|<x_c} dx\sqrt{-x^2 C_1 - a_* + \frac{a_*'^2}{2a_*''}}. \tag{42}$$

Integrating (42), one arrives at the imaginary potential for the back reacted gravity background

$$ImV_{Q\bar{Q}} = -\frac{1}{2\sqrt{2}\alpha'}(\frac{a_*'}{2a_*''} - \frac{a_*}{a_*'}), \tag{43}$$

with

$$
\begin{aligned}
a'_* &= r_*^4 f'(r_*) + 4r_*^3 f(r_*),\\
a''_* &= r_*^4 f''(r_*) + 8r_*^3 f'(r_*) + 12r_*^2 f(r_*),\\
f'(r_*) &= 8mr_*^{-5} + 2br_*^{-4},\\
f''(r_*) &= -40mr_*^{-6} - 8br_*^{-5},
\end{aligned}
\tag{44}
$$

where, for convenience, we set $R = 1$. Note that when b is put to zero, the result of $\mathcal{N} = 4$ SYM plasma [17] is reproduced.

4. Numerical results

Before numerical computation, we discuss the regime of applicability of this model. First, the imaginary potential should be negative, yielding

$$
\frac{a'_*}{2a''_*} - \frac{a_*}{a'_*} > 0, \tag{45}
$$

results in

$$
\varepsilon > \varepsilon_{min}, \tag{46}
$$

with $\varepsilon = r_h/r_*$, where the value of ε_{min} can be evaluated numerically.

The second limitation relates to the maximum value of LT (say LT_{max}). From (14) and (26), one has

$$
LT = \frac{6r_h^3 - b}{3\pi r_h^2} \int_{r_*}^{\infty} dr \sqrt{\frac{a(r_*)}{a^2(r) - a(r)a(r_*)}}. \tag{47}
$$

In Fig. 2, we plot LT as a function of ε for different values of b. From these figures, one can see that in each plot there exists a LT_{max}, corresponds to $\varepsilon = \varepsilon_{max}$. In fact, for $\varepsilon > \varepsilon_{max}$, there may be other string configurations (which are not solutions of the Nambu-Goto action) that contribute to the calculation of the Wilson loops besides the semiclassical U-shaped string configuration [39]. Here we are mostly interested in the case of $\varepsilon < \varepsilon_{max}$. Briefly, the domain of applicability of (43) is $\varepsilon_{min} < \varepsilon < \varepsilon_{max}$.

Now let's discuss results. First, we analyze how back reaction affects the inter-distance of $Q\bar{Q}$. From Fig. 2, one sees that increasing b leads to decreasing LT, implying the inclusion of back reaction reduces the inter-distance.

Next, we study how back reaction influences the imaginary potential. To that end, we plot $\mathrm{Im}V/(\sqrt{\lambda}T)$ versus LT for different values of b in Fig. 3. One can see that in each curve $\mathrm{Im}V/(\sqrt{\lambda}T)$ starts at a LT_{min}, corresponding to ε_{min}, and ends at a LT_{max}, corresponding to ε_{max}. Moreover, for increasing b the onset of $\mathrm{Im}V/(\sqrt{\lambda}T)$ happens at smaller LT. As discussed in [20], if the onset of the imaginary potential occurs for smaller LT, the suppression will be stronger. Thus, one concludes that the presence of back reaction enhances the quarkonia dissociation.

Also, one can analyze the effect of back reaction on the thermal width. Usually, a larger absolute value of imaginary potential corresponding to a larger thermal width [17,18]. Since the absolute value of $\mathrm{Im}V/(\sqrt{\lambda}T)$ decreases with b, one infers that the inclusion of back reaction decreases the thermal width.

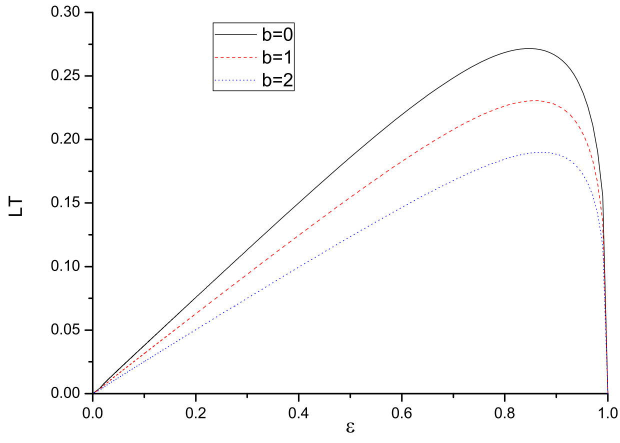

Fig. 2. LT versus ε. In the plots from top to bottom, $b = 0, 1, 2$, respectively.

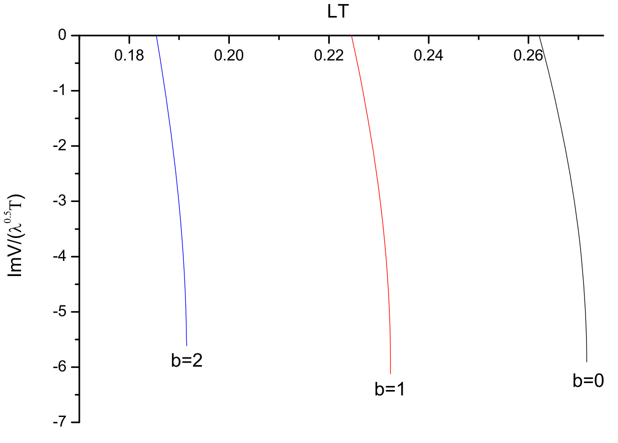

Fig. 3. $\mathrm{Im}V/(\sqrt{\lambda}T)$ versus LT. In the plots from right to left, $b = 0, 1, 2$, respectively.

5. Conclusion and discussion

In this paper, we studied the imaginary potential of heavy quarkonia in a back reacted gravity background, which is dual to a system of large number of heavy quarks uniformly distributed over the $\mathcal{N} = 4$ SYM plasma. The motivation rests on the earlier studies on the drag force [32] and jet quenching parameter [33] in such theories. It is observed that the presence of back reaction decreases the absolute value of the imaginary potential thus decreasing the thermal width. One step further, the inclusion of back reaction enhances the quarkonia dissociation.

However, several problems remain. The most pertinent how to study the back reaction with respect to the interaction between the heavy quarks in the cloud and the back reaction to the spacetime geometry? Moreover, how to study such effects in the systems that are not conformal? (since SYM plasma differs from QGP mainly due to the lack of a dynamical breaking of conformal symmetry). We hope to report our progress in this regard in the near future.

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