Jet quenching parameter of Sakai-Sugimoto model

Yi-hong Gao, Wei-shui Xu and Ding-fang Zeng

Institute of Theoretical Physics, Chinese Academy of Science

Abstract

Using gauge theory/string duality, we calculated the jet quenching parameter $\hat{q}$ of the Sakai-Sugimoto model in various phases. Being different from the $\mathcal{N}=4$ SYM theory where $\hat{q} \propto T^3$, we find that $\hat{q} \propto T^4/T_d$, where $T_d$ is the critical temperature of the confined/deconfined phase transition. By analyzing the $\hat{q}$ in different phases of this theory, we get better understanding about some statements in previous works, such as the non-universality and the explanation of discrepancies between the theory predictions and experiments.
1 Introduction

Experimental relativistic heavy ion collisions produced many evidences signalling the quark gluon plasma’s formation [1,2] at RHIC. One of the characteristic features discovered from data is the suppression of heavy quark’s spectrum in high-$p_T$ region [3,4]. This is the jet quenching phenomenon. Its cause is the medium induced heavy quarks’ energy loss when they move through QGP. So successful models explaining this phenomenon usually involve a medium sensitive “jet quenching parameter” $\hat{q}$. While in [5], the authors proposed a first principle, non-perturbative quantum field theoretic definition for it. Using their new definition and AdS/CFT correspondence [6,7,8], these authors calculated the jet quenching parameter of the thermal $\mathcal{N} = 4$ supersymmetric Yang-Mills theory. After some reasonable parameters’s choice adapted to the Au-Au collisions at RHIC and the universality assumption $\hat{q}_\text{QCD} \approx \hat{q}_\mathcal{N}=4$, they find that the theoretically predicted $\hat{q}$ is less than that suggested by experiments. And the author of [9] generalizes the proposal of [5] to the strongly coupled non-conformal gauge theory plasma and find that instead of universal, the jet quenching parameter is gauge theory specific. They explicitly show that the jet quenching parameter increases as one goes from a confining gauge theory to a conformal theory. So they concluded that the reasons about the jet quenching parameter predicted by [5] being less than the experimental measures is probably due to the existence of some extra energy loss mechanisms of the heavy quarks besides gluon radiation.

In order to investigate whether the jet quenching parameter is universal or not, the calculations of this parameter in some different gauge theories are meaningful and will help us to understand discrepancies between theory and experiments more deeply. So after [5] and [9], many other works related to the jet quenching parameter appear, see e.g. [10], [11], [12], [13], [14] and [15]. One common feature of these works is that, in the dual gravity description of the corresponding gauge field theory, the background metric usually involves an asymptotically $AdS_5$ component.

An interesting question is, how will the jet quenching parameter behave in a gauge theory whose dual gravity descriptions involves no $AdS_5$ component? In this paper, we will investigate this problem by using the Sakai-Sugimoto model [17]. It is a holographic model of the low energy QCD constructed from the intersecting branes. whose most striking feature is its phase structure [18]: (i) the zero and low temperature phase or confining phase in which all fundamental field degrees of freedom is confined; (ii) in the deconfined phase, there exists the chiral symmetry breaking phase and the chiral symmetry restoration phase.

\[1\] For more references, one can see the paper [16]
We will follow the same routine as in [5], and calculate the jet quenching parameter. After doing some analysis this parameter in the different phase of this model, we can get more deep understanding about the jet quenching phenomenon in the thermal QCD-like theory.

The organization of this paper is as follows, we will briefly review the Sakai-Sugimoto model in the next section, then calculate the jet-quenching parameter of this model in different phases in the section 3, while the last section contains the main conclusions of the paper.

2 Brief review of Sakai-Sugimoto model

In the [17], Sakai and Sugimoto constructed a holographic QCD-like model from the brane components D4, D8 and $\overline{D8}$ in type IIA string theory. In this brane construction, the low energy effective theory in the intersecting dimension is a $U(N_c)$ gauge theory with $N_f$ flavors. And there also exists a global $U(N_f)_L \times U(N_f)_R$ chiral symmetry.

In the strong coupling region, we can use the supergravity approximation to analyze the underlying physics. Thus, in the zero-temperature phase, we can use the D8 brane to probe the following background

$$ds_c^2 = H^{-\frac{1}{2}}(-dt^2 + d\vec{x} \cdot d\vec{x} + f dx_4^2) + H^\frac{1}{2}(f^{-1} du^2 + u^2 d\Omega_4^2),$$

$$H = \frac{R_{D4}^3}{u^3}, \quad f = 1 - \frac{u_\Lambda^3}{u^3}, \quad R_{D4}^3 = N_c \pi g_s l_s^3,$$

$$e^\phi = g_s \left[ \frac{u^3}{R_{D4}^3} \right]^{1/4}, \quad F(4) = \frac{2\pi N_c}{V_4} \epsilon_4,$$  \hspace{1cm} (1)

where $t$, $\vec{x}$ are four uncompactified world-volume coordinates of the D4 branes, and $x_4$ denotes the compactified world-volume coordinate. The sub-manifold spanned by $x_4$ and $u$ is singular as $u \rightarrow u_\Lambda$, to avoid this singularity, the periodical identification of $x_4$ should be

$$x_4 \sim x_4 + 2\pi R, \quad 2\pi R = \frac{4\pi}{3} \left[ \frac{R_{D4}^3}{u_\Lambda} \right]^{1/2}.$$ \hspace{1cm} (2)

While to the high temperature phase of the gauge theory i.e $T > (2\pi R)^{-1} \equiv T_d$ (where $T_d$ is the critical temperature of the Hawking-Page phase transition), the dual gravity system’s

\[\text{In [20, 19], the authors have studied the drag force and screening length problems in this holographic model.} \]
The background becomes

$$ds^2_8 = H^{-\frac{1}{2}}(-f dt^2 + d\vec{x} \cdot d\vec{x} + dx_4^2) + H^{\frac{1}{2}}(f^{-1}du^2 + u^2d\Omega^2_4)$$

$$H = \frac{R^3_{D4}}{u^3}, \quad f = 1 - \frac{u_T^3}{u^3}. \quad (3)$$

The others are same as in (1), and here the Hawking temperature is $T = \frac{3}{4\pi} \left[ \frac{u_T}{R_{D4}} \right]^{1/2}$.

As investigated in [18], there exists two phases in this deconfined phase. We can use the $D8$- and $\overline{D8}$-branes to probe the background (3). The world-volume coordinates of the $D8$- and $\overline{D8}$-branes consists of $t$, $\vec{x}$, $u$ and $\Omega_4$. Assuming that as $u \to \infty$ the D8 branes sit at $x_4 = 0$ while the $\overline{D8}$-branes at $x_4 = D$, then the induced metric on the probe $D8$- or $\overline{D8}$-branes reads

$$ds^2_8 = H^{-\frac{1}{2}}(-f dt^2 + d\vec{x} \cdot d\vec{x}) + (H^{-\frac{1}{2}}x_4^2 + H^{\frac{1}{2}}f^{-1})du^2 + H^{\frac{3}{2}}u^2d\Omega^2_4. \quad (4)$$

The DBI action of the probe $D8$-branes gives equation of motion,

$$\frac{u^4f}{\sqrt{f + H/[dx_4/du]^2}} = u_0^4\sqrt{f(u_0)}$$

$$\Rightarrow \frac{dx_4}{du} = \pm \left[ \frac{H}{f([u^8f]/[u_0^8f_0] - 1)} \right]^{1/2} \quad (5)$$

where $u_0$ is the $u$-coordinate of a point at which the $x_4'(u) = 0$. As the result,

$$x_4(u) = \frac{D}{2} \pm \left[ \frac{R^3_{D4}}{u_0} \right]^{1/2} \int_1^{\frac{u_T}{u_0}} dx \left[ \frac{x^{-3}}{(1 - y_T^3x^{-3})[x^8(1 - y_T^3) / (1 - y_T^3) - 1]} \right]^{1/2} \quad (6)$$

where $y_T = \frac{u_T}{u_0}$. In this case, the $D8$- and $\overline{D8}$ are connected smoothly at $x_4 = D/2$ point.

However, the configuration constrained by

$$\partial_u x_4 \equiv 0 \quad (7)$$

also satisfy eq(6). This means that the $D8$ and $\overline{D8}$ are located at fixed positions

$$x_4(u) \equiv 0 \quad \text{or} \quad x_4(u) \equiv D. \quad (8)$$

Between the phases (6) and (8), as pointed out by [18], there exists a first order chiral phase transition. Under a critical temperature $T_{\chi SB} \simeq 0.154/D$, the chiral symmetry is broken.
However, above this temperature, the symmetry will be restored. We depict the relevant profile in the left hand side of Figure 1.

![Figure 1: The chiral symmetry is broken in the left one, however, it is restored in the right one.](image)

### 3 The jet quenching parameter

#### 3.1 General definitions

By the definition of [5], the jet quenching parameter $\hat{q}$ is related to the expectation value of a light-like wilson loop $< W^A(C) >$ in the adjoint representation whose contour $C$ composed of a rectangle of length $L^-$ and width $L$ with $L << L^-$ and $L << T^{-1}$. Motivated by dipole approximation [20]

$$< W^A(C) > \approx \exp \left[ -\frac{1}{4\sqrt{2}} \hat{q} L^- L^2 \right].$$

(9)

The author of [5] defined $\hat{q}$ as the coefficient of the $L^- L^2/(4\sqrt{2})$ term in $\ln < W^A(C) >$. So

$$\hat{q} = \frac{-4\sqrt{2}}{L^- L^2} \ln < W^A(C) >.$$

(10)

The factor of $\sqrt{2}$ in the above equations comes from the light-cone coordinate definition $x^\pm = (t \pm x^1)/\sqrt{2}$. If a different definition $x^\pm = t \pm x^1$ are used, the factor of $\sqrt{2}$ in the equation [9] will not appear, see for e.g. [9].
So, let us consider such a Wilson loop following the same routine of [5]. In the planar limit, wilson loops in the adjoint representation is related to that in the fundamental representation through

\[ \ln < W^A(C) >= 2 \ln < W^F(C) >. \]  \hspace{1cm} (11)

According to the gauge theory/string duality, see for e.g. [21], [22], [23] and [24], \( < W^F(C) > \) can be given by

\[ < W^F(C) >= \exp[i\hat{S}(C)] \]  \hspace{1cm} (12)

where \( \hat{S} \) is related to the extremal action \( S \) of a fundamental string whose world-sheet on the boundary of the background space-time have \( C \) as the boundary, i.e.

\[ S = -\frac{1}{2\pi\alpha'} \int d^2 \sigma \sqrt{-\gamma}, \]
\[ \gamma = \det \gamma_{ab}, \quad \gamma_{ab} = \partial_a X^\mu \partial_b X^\nu G_{\mu\nu} \]  \hspace{1cm} (13)

which is just the Nambu-Goto action of the string. In the above equations, \( X^\mu \) is the embedding of the fundamental string in the background space-time with metric \( G_{\mu\nu} \). From gauge theory, since the action \( S \) contains the self-energy \( S_{sef} \) of quarks attached on the two ends of the string, it is usually divergent. In order to describe the interactions between the two quarks, we need subtract the self-energy part. The final part \( \hat{S} \) will be

\[ \hat{S} = S - S_{sef} \]  \hspace{1cm} (14)

### 3.2 String configuration

According to the general routines in the last section, in order to calculate the jet quenching parameters of the Sakai-Sugimoto model in different phases, we can write the background space-time (1) and (3) in the light-cone coordinates. The zero temperature background becomes

\[ ds^2 = H^{-\frac{1}{2}}(-2dx^+dx^- + dx_1^2 + f dx_3^2) + H^{\frac{1}{2}}(f^{-1}du^2 + u^2d\Omega_4^2) \]  \hspace{1cm} (15)
and the high temperature one is

\[ ds^2 = H^{-\frac{1}{4}} \left( -\frac{1+f}{2} dx^+ dx^- + \frac{1-f}{4} [dx^2 + dx^{-2}] + dx_\perp^2 + dx_4^2 \right) \\
+ H^{\frac{1}{4}} \left( f^{-1} du^2 + u^2 d\Omega^2_4 \right). \quad (16) \]

To find the embedding of strings into the above backgrounds, we give the following ansatz for the string configuration,

\[ X^- = \tau, \; X^2 = \sigma, \]
\[ u(\sigma), \; X^4(\sigma) \quad (17) \]

Under this ansatz, the Nambu-Goto action of the string vanishes in the zero temperature metric (15), and in the high temperature background (16) reads

\[ S = -\frac{1}{2\pi\alpha'} \int d^2\sigma \sqrt{\frac{f-1}{2} \left( \frac{1}{H} + \frac{u^2}{f} + \frac{u^2}{H} \left[ \frac{dx_4}{du} \right]^2 \right)} \quad (18) \]

where \( u' \) denotes \( \frac{du}{d\sigma} \).

By minimizing the Nambu-Goto action, we can find the equation of motion and then the solution of the required string configuration. Obviously, since the action equals to zero identically in the zero-temperature phase, the corresponding string configuration satisfying required boundary conditions cannot be determined uniquely. Notice that it is \( \hat{S} \) instead of \( S \) that determines the jet quenching parameter. Hence this means the jet quenching parameter \( \hat{q} \) will be equal to zero. Since the jet quenching parameter is designed to describe the property of medium through which the test heavy quark moves. The vanishing jet quenching parameter \( \hat{q} \) means that the medium will look transparent to the test heavy quarks in the confining phase.

For the high temperature phase, we can get the equation of motion from the action (18) which reads

\[ u'^2 = \frac{c^2 f}{H + (dx_4/du)^2 f} \quad (19) \]

In deriving these equations, we used the fact that Lagrangian \( \mathcal{L} \equiv \sqrt{-\gamma} \) in the equation (18) does not contain \( \sigma \) explicitly. So the effective “hamiltonian” \( \mathcal{H} \equiv u' \frac{\partial \mathcal{L}}{\partial u'} - \mathcal{L} \) is a conserved
quantity. The integral constant $c$ characterizes the profile of string. If the $dx_4/du = 0$, then the equation (19) will reduce to

$$u'^2 = \frac{c^2f}{H}$$

(20)
due to the fact that $dx_4/du = 0$. In this case, the test string lies in the quark gluon plasma with chiral symmetry restored.

We require the string has the profile depicted in figure 2 intuitively. The end points of this string are located on the $D8$-branes world-volume. And its projection on this worldvolume will form a light-like Wilson loop. It goes from the world volume of the $D8$-branes at $u = \infty$, towards the black hole horizon then turn around at some lowest point and come back to the $D8$-branes world-volume again, i.e. the string has a U-shape in the $x_2-u$ plane.

From equation (20), we find that in the chiral symmetry restored phase the only zero point occurs at the black hole horizon. It means the turning point of the string will exactly reach the horizon. And in this phase, the length of the string is

$$L = 2 \int \frac{du}{u'} = \frac{u_T}{c} \left[ \frac{R_{D8}^3}{u_T^3} \right]^{1/2} \frac{2\sqrt{\pi} \Gamma[1/6]}{3 \Gamma[2/3]}$$

(21)

where we have used the integral

$$\int_1^\infty \frac{dx}{\sqrt{x^3 - 1}} = \frac{\sqrt{\pi} \Gamma[1/6]}{3 \Gamma[2/3]}.$$
While in the broken chiral symmetry phase, equation (19) has two zero points, one at the black hole horizon \( u = u_T \), the other at \( u = u_0 > u_T \). To assure the positivity of \( u'' \), the lowest point of the string can only reach \( u = u_0 \). Then, the length \( L' \) of string is

\[
L' = \frac{2}{c} \int \frac{\left[ H + (dx_4/du)^2 f \right]}{f} \cdot A[y_T]^{1/2} = \frac{u_T}{c} \left[ \frac{R_{D4}^3}{u_T^3} \right]^{1/2} \frac{2\sqrt{\pi} \Gamma[1/6]}{3 \Gamma[2/3]} \cdot A[y_T]
\]

with the definition

\[
\int_{u_0/u_T}^{\infty} \frac{dx}{x^3 - 1 - u_0^8 f_0/u_T^5 x^{-5}} = A[y_T] \cdot \frac{\sqrt{\pi} \Gamma[1/6]}{3 \Gamma[2/3]}
\]

where the \( y_T = u_T/u_0 \). Obviously, if \( A[y_T] = 1 \), then the length \( L' \) of the string in the broken chiral symmetry phase will reduce to the corresponding one in the symmetry restored phase. The numerical behavior of \( A[y_T] \) is displayed in the left hand side of Figure 3. From the figure we can see the behavior of the function \( A[y_T] \) with the variable \( y_T \). When \( y_T = 1 \), \( A[y_T] \) will reach the maximal value 1. This means that at this point the chiral symmetry broken phase will turn into the chiral symmetry restored phase.

Figure 3: The left panel, the numerical behavior of \( A[y_T] \) defined in the equation (24). The right panel, the distance \( D \) between the \( D8- \) and \( \overline{D8}\)-branes v.s. \( y_T \). The distance \( D \) is measured by the temperature inverse, \( T^{-1} \). When \( y_T << 1 \), \( D \propto y_T^{1/2} \) which can be looked out from the equation (6).
3.3 The jet quenching parameter in various phases

Now we can use the routines of subsection 3.1 to calculate the jet quenching parameters of the Sakai-Sugimoto model. We start with the Nambu-Goto action,

\[ S = \frac{L - \cdot L'}{2\sqrt{2\pi\alpha'}} \left[ \frac{u_T^3}{R_{D4}^3} \right]^{1/2} \cdot \sqrt{1 + c^2} \]  \hspace{1cm} (25)

and the self-energy of the string is

\[ S_{sef} = \frac{2}{2\pi\alpha'} \int d\tau du \sqrt{-G_{uu}} \]  
\[ = \frac{L - u_T}{\sqrt{2\pi\alpha'}} \sqrt{\frac{\pi}{3}} \Gamma^{[1/6]} \Gamma^{[2/3]} \cdot A[y_T] \]  \hspace{1cm} (26)

Then according to the equation (10), the jet quenching parameter can be expressed as

\[ \hat{q} = 2(S - S_{sef}) \cdot \frac{4\sqrt{2}}{L - L'^2} = \frac{4}{\pi\alpha'} \left( \left[ \frac{u_T^3}{R_{D4}^3} \right]^{1/2} \frac{(1 + c^2)^{1/2}}{L'} - \frac{u_T}{L'^2} \frac{2\sqrt{\pi} \Gamma^{[1/6]} \Gamma^{[2/3]} \cdot A[y_T]}{3} \right) \]  \hspace{1cm} (27)

The right hand side of this equation involves the undetermined integration constant \( c \). But it can be eliminated by the relations expressed in the equation (23). And since the wilson loop satisfies the condition \( L'T \to 0 \), we find

\[ \hat{q} = \frac{64\sqrt{\pi^3} \Gamma^{[2/3]} \Gamma^{[1/6]} A[y_T]}{27} \frac{\lambda}{T^4/T_d} \]  \hspace{1cm} (28)

where we have used the relation \( \lambda = g_{Y_M}^2 N_c, g_{Y_M}^2 = 2\pi g_s l_s / R \), and the critical temperature \( T_d = 1/(2\pi R) \). If choosing the parameter \( A[y_T] = 1 \), then the above results will reduce to the corresponding ones in the chiral symmetry restoration phase. However, if \( A[y_T] \neq 1 \), all the results in the above are in the chiral symmetry broken phase.

Let us make some comments on these results in the following. First, we compare the jet quenching parameters in the chiral symmetry broken phase and that in the chiral symmetry restored one. Since \( A[y_T] < 1 \), we know that the jet quenching parameter \( \hat{q} \) in the chiral symmetry broken phase is larger than the one in the chiral symmetry restored phase. As is known, in the chiral symmetry broken phase, the gauge system is a mixture of quark-gluon plasma and hadrons; while in the chiral symmetry restored phase, the gauge system is a pure quark-gluon plasma with all hadrons resolved. So if the dominant mechanism of heavy quarks energy loss is due to gluon radiation induced by the medium, we expect the jet quenching
parameter in the chiral symmetry broken phase be less than that in the chiral symmetry restored one, since the gluons density in the former is less than that in the latter. So we think that our result about the parameter $\hat{q}$ being more large in the chiral symmetry broken phase may be looked as an evidence supporting the viewpoint of [9] which states that, the reasons of the jet quenching parameter found in [5] being less than the experimental measures are due to existence of some extra energy loss mechanisms (such as binary collisions) of the heavy quarks as they move through the mediums. [28] even suggest that the dominant mechanism is not gluon radiation but binary collisions.

Second, let us make a comparison between the jet quenching parameter of Sakai-Sugimoto model plasmas with the $N = 4$ SYM theory plasmas. The most striking difference is their temperature dependence. The former, $\hat{q}_{SS} \propto T^4$ at a fixed $T_d$ while the latter, $\hat{q}_{N=4SYM} \propto T^3$. This is an evidence that the jet quenching parameter in the different gauge theories is not universal. If experiments can measure the jet quenching parameter v.s. temperature relation explicitly, this may be a good starting point for further exploration of the meaning of the jet quenching parameter.

Finally, we make a little numerics to see whether the Sakai-Sugimoto theory plasma fit experiments more better or even worse than $N = 4$ SYM plasmas. For this purpose, we rewrite the equation (28) as

$$\hat{q} = 3.21 \frac{A[y_T]}{\lambda} \frac{T^4}{T_d}$$

(29)

If letting $A[y_T] = 1$, it reduces to

$$\hat{q} = 3.21 \lambda T^4 / T_d$$

(30)

in the chiral symmetry restoration phase.

Taking $N_c = 3$, $g^2_{YM} = 0.5$ and $T_d = 300$ Mev, we find, in order to obtain $\hat{q} = 5 \text{GeV}^2/\text{fm}$, the corresponding temperature of the quark gluon plasma in the chiral symmetry restoration phase is 280 Mev, which is smaller than expected [25]. Equivalently, the jet quenching parameter on the basis of these parameters are less than that suggested by RHIC data, note that, $T \sim \left[ \frac{\hat{q}_T}{N_c g^2_{YM}} \right]^{1/4}$. However, in the chiral symmetry broken phase, the corresponding temperature is $280 \cdot A^4[y_T]$ Mev. As long as we choose a appropriate $A[y_T]$, as a result, the jet quenching parameter can be equal to the one suggested by RHIC data.

To understand what the smallness of $A[y_T]$ means. Let us turn back to the Sakai-Sugimoto model for a while. From figure [8] we know that smaller $A[y_T]$ corresponds to smaller $y_T$. From
the equation (6) we know that \( y_T \) is related with the distance \( D \) between the \( D8 \) and \( \overline{D8} \)-branes. We numerically integrated this equation and plot the \( D \) v.s. \( y_T \) relation in the right panel of Figure 3. From the figure we see that in the small \( y_T \) region, \( D \) is an increasing function of \( y_T \). So the smallness of \( y_T \) means the small distance between the \( D8 \)- and \( \overline{D8} \)-branes. Hence from aspects of the microscopic construction of Sakai-Sugimoto model, to enhancing the jet quenching parameter, we only need to make the \( D8 \)- and \( \overline{D8} \)-branes sit as near as possible.

4 Summary

In this paper, we calculated the jet quenching parameters of the Sakai-Sugimoto model in various phases. In this holographic model, the gravity metric is different from the asymptotic \( AdS_5 \) spacetimes. After some analysis, we find that the jet quenching parameter \( \hat{q} \) in this holographic model is proportional to the fourth power of the temperature at a fixed confining/deconfining temperature \( T_d, \hat{q} \propto T^4/T_d \). By comparing the jet quenching parameter of this model in different phases with that of the \( N = 4 \) SYM plasmas, we get more deep understanding about the following statements in previous works, (i) the jet quenching parameters in different gauge theories are not universal, bur are specific for a given gauge theory; (ii) the discrepancies about the jet quenching parameter between theoretically calculating and experimental data is probably due to some extra energy loss mechanisms except gluon radiations of the heavy quarks when they move through the quark gluon plasmas.

References

[1] J. Adams(Star Collaboration) et. al., “Experimental and theoretical challenges in the search for the quark gluon plasma: The STAR collaboration’s critical assessment of the evidence from RHIC collisions”, Nucl. Phys. A757 (2005) 102–183, nucl-ex/0501009.

[2] K. Adcox et. al., “Formation of dense partonic matter in relativistic nucleus nucleus collisions at RHIC: Experimental evaluation by the PHENIX collaboration”, Nucl. Phys. A757 (2005) 184–283, nucl-ex/0410003.

[3] S. S. Adler et. al., “Modifications to di-jet hadron pair correlations in \( \text{Au + Au} \) collisions at \( s(\text{NN})^{(1/2)} = 200-\text{GeV} \)”, Phys. Rev. Lett 96 (2006) 032301, nucl-ex/0510047.
[4] J. Bielcik(Star Collaboration) et. al., “Centrality dependence of heavy flavor production from single electron measurement in $s(\text{NN})^{(1/2)} = 200$-GeV Au + Au collisions”, Nucl. Phys. A774 (2006) 697, [nucl-ex/0511005]

[5] H. Liu, K. Rajagopal, and U. A. Wiedemann, “Calculating the jet quenching parameter from AdS/CFT”, [hep-ph/0605178]

[6] J. Maldacena, “The Large N Limit Of Superconformal Field Theories And Supergravity”, [hep-th/9711200]

[7] E. Witten, “Anti-de Sitter space, thermal phase transition and confinement in gauge theories”, Adv.Theor.Math.Phys.2(1998)505, [hep-th/9803131]

[8] O. Aharony, S. S. Gubser, J. M. Maldacena, H. Ooguri and Y. Oz, “Large-N Field Theories, String Theory and Gravity”, Phys. Rep. 323(2000)183, [hep-th/9905111]

[9] A. Buchel, “On jet quenching parameters in strongly coupled non-conformal gauge theories”, Phys. Rev. D74(2006)046006, [hep-th/0605178]

[10] J. F. Vazquez-Poritz, “Enhancing the jet quenching parameter from marginal deformations”, [hep-th/0605296]

[11] E. Cáceres and A. Güijosa, “On Drag Forces and Jet Quenching Strong-Coupled Plasmas”, [hep-th/0606134]

[12] F. L. Lin, T. Matsuo, “Jet Quenching Parameter in Medium with Chemical Potential from AdS/CFT”, [hep-th/0606136]

[13] S. D. Avramis and K. Sfetsos, “Supergravity and the jet quenching parameter in the presence of R-charge densities”, [hep-th/0606190]

[14] N. Armesto, J. D. Edelstein and J. Mas, “Jet quenching at finite ’t Hooft coupling and chemical potential from AdS/CFT”, JHEP 0609:039,2006, [hep-ph/0606245]

[15] E. Nakano, S. Teraguchi and Wen-Yu Wen, “Drag force, jet quenching, and AdS/QCD”, [hep-ph/0608274]

[16] H. Liu, K. Rajagopal and U. A. Wiedemann, “Wilson loops in heavy ion collisions and their calculation in AdS/CFT,” JHEP 0703, 066 (2007) [arXiv:hep-ph/0612168].
[17] T. Sakai and S. Sugimoto, “Low Energy Hadron Physics in Holographic QCD”, Prog. Theor. phys. 113(2005)843-882, [hep-th/0412141].

[18] O. Aharony, J. Sonnenschein and S. Yankielowicz, “A Holographic Model of Deconfinement and Chiral Symmetry Restoration”, [hep-th/0604161]. A. Parnachev and D. A. Sahakyan, “Chiral phase transition from string theory,” Phys. Rev. Lett. 97, 111601 (2006) [arXiv:hep-th/0604173].

[19] K. Peeters, J. Sonnenschein and M. Zamakar, “Holographic melting and related properties of mesons in a quark gluon plasma”, [hep-th/0606195].

[20] B. G. Zakharov, “Radiative energy loss of high-energy quarks in finite size nuclear matter and quark-gluon plasma”. JETP Lett 65 (1997) 615, [hep-ph/9704255].

[21] J. M. Maldacena, “Wilson loops in large N field theories,” Phys. Rev. Lett. 80 (1998) 4859–4862, [hep-th/9803002].

[22] S.-J. Rey and J.-T. Yee, “Macroscopic strings as heavy quarks in large N gauge theory and anti-de Sitter supergravity”, Eur. Phys. J. C22(2001) 379–394, [hep-th/9803001].

[23] S.-J. Rey, S. Theisen and J.-T. Yee, “Wilson-Polyakov loop at finite temperature in large N gauge theory and anti-de Sitter supergravity”, Nucl. Phys. B527 (1998) 171–186, [hep-th/9803135].

[24] A. Brandhuber, N. Itzhaki, J. Sonnenschein, and S. Yankielowicz, “Wilson loops in the large N limit at finite temperature”, Phys. Lett. B434(1998) 36–40, [hep-th/9803137].

[25] P. F. Kolb and U. Heinz, “Hydrodynamic description of ultrarelativistic heavy-ion collisions”, nucl-th/0305084.

[26] P. Talavera, “Drag force in a string model dual to large-N QCD”, [hep-th/0610179].

[27] P. C. Argyres, M. Edalati, J. F. Vazquez-Poritz, “No-Drag String Configurations for Steadily Moving Quark-Antiquark Pairs in a Thermal Bath”, [hep-th/0608118].

[28] H. van Hees, V. Greco and R. Rapp, “Heavy-Quark Probes of the Quark-Gluon Plasma at RHIC”, Phys.Rev. C73 (2006) 034913, [nucl-th/0508055].