Heavy quark potential from deformed $AdS_5$ models

Zi-qiang Zhang, De-fu Hou, and Gang Chen

1 School of mathematics and physics, China University of Geosciences (Wuhan), Wuhan 430074, China
2 Key Laboratory of Quark and Lepton Physics (MOE), Central China Normal University, Wuhan 430079, China

In this paper, we investigate the heavy quark potential in some holographic QCD models. The calculation relies on a modified renormalization scheme mentioned in a previous work of Albacete et al. After studying the heavy quark potential in Pirner-Galow model and Andreev-Zakharov model, we extend the discussion to a general deformed $AdS_5$ case. It is shown that the obtained potential is negative definite for all quark-antiquark separations, differs from that using the usual renormalization scheme.

PACS numbers: 12.38.Lg, 12.39.Pn, 11.25.Tq

I. INTRODUCTION

Heavy quark potential is an important quantity in the study of Quark Gluon Plasma (QGP) in heavy ion collisions carried out at RHIC or LHC, since the melting of heavy mesons is usually considered to be one of the signatures for QGP formation. However, much experiment data indicates that QGP is strongly coupled [1–3]. Thus, calculational tools for strong coupled, real-time QCD dynamics are required. Such tools are now available via the AdS/CFT correspondence [4–6].

AdS/CFT, the duality between the type IIB superstring theory formulated on $AdS_5 \times S^5$ and $\mathcal{N}=4$ SYM in four dimensions, has yielded many important insights into the dynamics of strongly-coupled gauge theories. By using the AdS/CFT, Maldacena has calculated the heavy quark potential in vacuum for $\mathcal{N}=4$ SYM firstly [7]. There, it is found that for the $AdS_5$ space the energy shows a purely Coulombian (non confining) behavior, agrees with a conformal gauge theory. This proposal has attracted lots of interest. Soon after [7], there are many attempts to addressing heavy quark potential in this direction. For instance, the potential at finite temperature has been discussed in [8, 9]. The potential in different spaces is investigated in [10]. The subleading order corrections to this quantity are considered in [11, 12]. This quantity is also analyzed with some AdS/QCD models [13–16]. Other important results can be found, for example, in [17–19].

Interestingly, the heavy quark potential can also be obtained by using an alternative renormalization scheme [20]. There, instead of subtracting the self energy of two static quarks, it subtracts the real part of the action of the two quarks that separated by an infinite distance. It was argued [20] that for AdS/CFT the obtained potential of the two renormalization schemes is the same at zero temperature but differs at finite temperature. For finite temperature case, they find a smooth non-zero (negative definite) potential without a kink, and the potential has a non-zero imaginary (absorptive) part for separations $r > r_c = 0.870/\pi T$. This result is obviously different from that using the usual renormalization procedure [8, 9], where the resulting potential goes to zero above $r > r^*_c = 0.754/\pi T$.

As the alternative renormalization is very interesting and the heavy quark potential based on it has been discussed in the case of AdS/CFT [20], it is also of interest to extend the studies of [20] to the case of AdS/QCD. In this paper, we would like to see whether we can get a new and interesting result by using this alternative renormalization for AdS/QCD. In addition, we will compare the results with those from the usual renormalization scheme and explore the reasonability of the results by comparing with the experimental data. These are the main motivations of the present work.

The paper is organized as follows. In the next section, some holographic QCD models are briefly reviewed. In section 3, using the modified renormalization procedure, we calculate the heavy quark potential in Pirner-Galow model, Andreev-Zakharov model and a general deformed $AdS_5$ model one by one. The summary and discussion is given in section 4.

*Electronic address: zhangzq@cug.edu.cn
†Electronic address: houdf@mail.ccnu.edu.cn
‡Electronic address: chengang1@cug.edu.cn
II. DEFORMED ADS BACKGROUNDS

In this section, we briefly review some deformed AdS backgrounds in [22]. There are two main approaches in searching for the string description of realistic QCD: top-down approach, i.e. by deriving holographic QCD from string theory [22–26], and bottom-up approach, i.e. by examining holographic QCD from experimental data and lattice results [16, 27–31].

In the bottom-up approach, the most economic way is to use a deformed AdS$_5$ metric, which can describe the known experimental data and lattice results. The simplest holographic QCD model is the hard-wall AdS$_5$ model [35], which has a good description of the lightest meson spectra comparing with the experimental data, but cannot produce the Regge behavior for higher excitations. The soft-wall model [28] suggests that a negative quadratic dilaton term $-z^2$ in the action can produce the right linear Regge behavior of $\rho$ mesons or the linear confinement. To produce linear behavior of heavy flavor potential, a positive quadratic term modification [27] is suggested to be added in the deformed warp factor of the metric in Andreev-Zakharov model [21]. To consider the QCD running coupling effect into the modified metric, see in [16, 31]. To take into account the QCD $\beta$ function, see in [22].

III. HOLOGRAPHIC HEAVY QUARK POTENTIAL

Before clarifying the definition of the heavy quark potential, we briefly introduce the Polyakov loop for SU($N_c$) gauge theory. At the spatial location $\vec{r}$, the Polyakov loop is

$$L(\vec{r}) = \frac{1}{N_c} Tr[P e^{i \int_0^\beta d\tau A_4(\vec{r}, \tau)}],$$

(1)

where $\tau$ is the Euclidean time, $T$ is the temperature and $\beta = 1/T$.

The singlet $V_1(r)$ and adjoint $V_{adj}(r)$ in Euclidean space can be extracted from the connected correlator of two Polyakov loops [32–33]

$$\langle L(0)L^\dagger(\vec{r}) \rangle_c \propto e^{-\beta V_1(\vec{r})} + (N_c^2 - 1)e^{-\beta V_{adj}(\vec{r})}.$$  

(2)

To obtain the Polyakov loop correlator in AdS space, one needs to connect the open strings to the positions of Polyakov loops at the boundary of the AdS space in all possible ways [7–9]. There are two relevant configurations, hanging string and straight string, which can give two different saddle points of the Nambu-Goto action, $S_{NG}^{\text{hanging}}$ and $S_{NG}^{\text{straight}}$. In the large $N_c$ and large $\lambda$ limit, we have

$$\langle L(0)L^\dagger(\vec{r}) \rangle_c \propto e^{-S_{NG}^{\text{hanging}}} + (N_c^2 - 1)e^{-S_{NG}^{\text{straight}}}.$$  

(3)

Note that $S_{NG}^{\text{straight}}$ gives $V_{adj}(r)$ and $S_{NG}^{\text{hanging}}$ gives $V_1(r)$. In lattice simulations, $V_1(r)$ is regarded as the heavy-quark potential at finite temperature [34].

On the other hand, $V_1(r)$ can be extracted from the expectation value of the static Wilson loop operator between a static pair of $Q\bar{Q}$

$$e^{-iT V_1(r)} \sim \lim_{T \to \infty} W(C),$$  

(4)

with

$$W(C) = \frac{1}{N} Tr P e^{i \oint_C dx^\mu A_\mu},$$  

(5)

where $C$ is a closed loop in a 4-dimensional spacetime, which is usually taken as a rectangular loop of spatial and temporal extensions $r$ and $T$, respectively. $A_\mu$ is the non-Abelian gauge potential operator. $P$ enforces the path ordering along the loop $C$.

As a result, the heavy quark potential $V_1(r)$ (henceforth we call it $V_{QQ}$) can be written as

$$V_{QQ} = \frac{S_c}{T}.$$  

(6)

where $S_c$ is a regularized action associated with $S_{NG}^{\text{hanging}}$. 
A. In Pirner-Galow model

The background metric of the Pirner-Galow model can be written as \[ds^2 = h(z)\frac{R^2}{z^2}(-dt^2 + dz^2 + d\vec{x}^2),\] (7)

where \[h(z) = \frac{\log(\frac{1}{\epsilon})}{\log(1 + \Lambda z^2 + \epsilon)},\] (8)

with the IR singularity at \[z_{IR} = \sqrt{1 - \epsilon \Lambda^2},\] (9)

where \(\Lambda\) is related to the \(AdS_5\) radius as \(\Lambda = 1/R\), \(\epsilon\) denotes a dimensionless parameter, that is \(\epsilon \equiv l_s^2 \Lambda^2\), (10)

the parameter \(l_s\) can guarantee the conformal limit \(h(z) \to 1\) at \(z \to 0\).

The Pirner-Galow model has some properties: (1). The functional form of the warping factor (8) can coincide with the functional form of the QCD running coupling. (2). The new parameter \(\epsilon\) can be fitted by reproducing the short and long range Cornell potential. Due to this characteristic, evaluation of the heavy quark potential in this background can be considered as a good test of \(AdS/QCD\).

We now study the heavy quark potential in Pirner-Galow model using the metric of (7). The string action can reduce to the Nambu-Goto action

\[S = -\frac{1}{2\pi\alpha'} \int d\tau d\sigma \sqrt{-detg_{\alpha\beta}},\] (11)

with \[g_{\alpha\beta} = G_{\mu\nu} \frac{\partial X^\mu}{\partial \sigma^\alpha} \frac{\partial X^\nu}{\partial \sigma^\beta},\] (12)

where \(X^\mu\) and \(G_{\mu\nu}\) are the target space coordinates and the metric respectively, and \(\sigma^\alpha\) with \(\alpha = 0,1\) parameterize the world sheet. The string tension \(\frac{1}{2\pi\alpha'}\) is related to the ’t Hooft coupling constant by

\[\frac{R^2}{\alpha'} = \sqrt{\lambda}.\] (13)

By using the static gauge

\[t = \tau, \quad \sigma = x,\] (14)

and supposing the radical for the radial direction

\[z = z(x),\] (15)

the metric (7) can reduce to the following form:

\[ds^2 = \frac{R^2}{z^2} h(z)[-dt^2 + (1 + \dot{z}^2)dx^2].\] (16)

Then the Euclidean version of Nambu-Goto action in (11) is given by

\[S = \frac{T R^2}{2\pi\alpha'} \int dz \frac{h(z)}{z^2} \sqrt{1 + \dot{z}^2},\] (17)

where \(\dot{z} = \frac{dz}{dx}\). Now we identify the Lagrangian as

\[\mathcal{L} = \frac{h(z)}{z^2} \sqrt{1 + \dot{z}^2}.\] (18)
Note that $\mathcal{L}$ does not depend on $x$ explicitly, so we have the conserved quantity,

$$\mathcal{L} - \frac{\partial \mathcal{L}}{\partial \dot{z}} \dot{z}. \quad (19)$$

The boundary condition at $x = 0$ is:

$$\dot{z} = 0, \quad z = z_0, \quad (20)$$

which yields

$$\frac{h(z)}{z^2\sqrt{1 + \dot{z}^2}} = \frac{h(z_0)}{z_0^2}, \quad (21)$$

then a differential equation is derived,

$$\dot{z} = \frac{dz}{dx} = \sqrt{\left(\frac{h(z)z_0^2}{h(z_0)z^2}\right)^2 - 1}. \quad (22)$$

By integrating (22) the separation length $L$ of the inter-quark is given by

$$L = 2 \int_0^{z_0} \frac{dz}{\sqrt{\left(\frac{h(z)z_0^2}{h(z_0)z^2}\right)^2 - 1}}. \quad (23)$$

Plugging (22) into (19), the action for the heavy quark pair is obtained

$$S(L, T) = \frac{T R^2}{\pi \alpha'} \int_0^{z_0} dz \frac{h(z)\sqrt{1 + \dot{z}^2}}{z^2 \dot{z}}. \quad (24)$$

Usually, the regularization action is related to the subtraction of the self energy of the two quarks, but we now use the following subtraction

$$V(L) = \frac{S(L, T) - Re[S(L = \infty, T)]}{T}. \quad (25)$$

To compare with the case in [16], we choose $\epsilon = 0.48$, $\Lambda = 0.264 GeV$. It was argued [16] that $z_r \rightarrow 1.85 GeV^{-1}$ when $L \rightarrow \infty$. In other words, the potential contains no imaginary part in the range $0 < z_0 < z_r$. Then, from (24) and (25), the heavy quark potential in Pirner-Galow model can be written as

$$V = \frac{R^2}{\pi \alpha'} \int_0^{z_0} dz \frac{h(z)\sqrt{1 + \dot{z}^2}}{z^2 \dot{z}} - \frac{R^2}{\pi \alpha'} \int_0^{z_r} dz \frac{h(z)\sqrt{1 + \dot{z}_1^2}}{z^2 \dot{z}_1}, \quad 0 \leq z_0 \leq z_r, \quad (26)$$

where

$$\dot{z}_1 = \sqrt{\left(\frac{h(z)z_0^2}{h(z_r)z^2}\right)^2 - 1}. \quad (27)$$

To proceed further, we need to turn to numerical methods. Here we plot the heavy quark potential $V(L)$ against the inter distance $L$ in Fig.1. From the figure we can see clearly that the potential is negative for all quark-antiquark separations, going to zero at an infinite distance. If one uses the usual renormalization procedure, the resulting potential goes to zero above a certain critical separation and also has a kink [16].

**B. In Andreev-Zakharov model**

Next, we consider the heavy quark potential in Andreev-Zakharov model. The corresponding background metric is given by [21]

$$ds^2 = \frac{R^2}{z^2} h(z)(-f(z)dt^2 + d\vec{x}^2 + \frac{1}{f(z)}dz^2), \quad (28)$$
FIG. 1: Typical graph of $V(L)$. Here $\epsilon = 0.48$, $\Lambda = 0.264 \text{GeV}$.

with

$$h(z) = e^{\frac{cz^4}{2}}, \quad f(z) = 1 - \frac{z^4}{z_h^4}, \quad (29)$$

where $c$ is a deformation parameter. The horizon is located at $z = z_h$ with $z_0 < z_h$, where $z_0$ is a turning point or deepest position of the string in the bulk. The temperature of the black hole is given by

$$T = \frac{1}{\pi z_h}. \quad (30)$$

The Andreev-Zakharov model has some properties: (1). In this background linearized five-dimensional Yang-Mills equations are effectively reduced to Laguerre differential equation, which can help to fix the value of $c$ from the $\rho$ meson trajectory $^{27}$, therefore the metric contains no free parameter. (2). The positive quadratic term modification in the deformed warp factor $h(z)$ can produce linear behavior of heavy flavor potential from the usual renormalization scheme.

We now explore the heavy quark potential in this metric $(28)$ with the alternative renormalization scheme. Parallel to the case of the previous Pirner-Galow model, we call again the distance between quark anti-quark pair and the action for the quark pair as $L$ and $S(L, T)$, respectively. One finds

$$L = 2 \int_0^{z_0} \frac{dz}{\sqrt{\frac{z^4}{2} e^{cz^4 - z_0^4} \left(1 - \frac{z^4}{z_h^4}\right) - \left(1 - \frac{z^4}{z_h^4}\right)}}. \quad (31)$$

and

$$S(L, T) = \frac{TR^2}{\pi \alpha'} \int_0^{z_0} dz \frac{h(z) \sqrt{f + \dot{z}^2}}{z^2 \dot{z}}. \quad (32)$$

Likewise, $z_0$ also has an upper bound $z_{\text{max}}$, and the potential does not contain an imaginary part. It has been verified in $^{13}$ that $z_{\text{max}}$ can be related to $z_h$ as

$$z_{\text{max}} = z_h \sqrt{\frac{2}{3} \sin \left(\frac{1}{3} \arcsin \frac{T^2}{T_1^2}\right)}, \quad (33)$$

with

$$T_1 = \frac{1}{\pi} \sqrt{\frac{c}{\sqrt{27}}}. \quad (34)$$
Then, from (25) and (32) the heavy quark potential in Andreev-Zakharov model can be written as

$$V = \frac{R^2}{\pi \alpha'} \int_0^{z_0} dz \frac{h(z) \sqrt{f + \dot{z}^2}}{z^2 \dot{z}} - \frac{R^2}{\pi \alpha'} \int_0^{z_{\text{max}}} dz \frac{h(z) \sqrt{f + \dot{z}_1^2}}{z^2 \dot{z}_1}, \quad 0 \leq z_0 \leq z_{\text{max}},$$

(35)

where

$$\dot{z}_1 = \left[ \frac{4}{z^4} e^{c(z^* - z_{\text{max}}^*)} \right] \left( 1 - \frac{z^4_{\text{max}}}{z^4} \right) - \left( 1 - \frac{z^4}{z^4_{\text{max}}} \right).$$

(36)

After taking $c = 0.9 \text{GeV}^2$ [27], we can plot the heavy quark potential versus the separate distance of the quark anti-quark numerically. From Fig.2 one finds the obtained potential is also negative definite, differs from that using the usual renormalization procedure [13].

C. In a general deformed $AdS_5$ model

To proceed, we extend the discussion to a general deformed $AdS_5$ model. The metric is given by

$$ds^2 = \frac{R^2}{z^2} h(z)(-f(z)dt^2 + d\vec{x}^2 + \frac{1}{f(z)} dz^2),$$

(37)

where the deformed warp factor $h(z)$ and the function $f(z)$ should satisfy the following conditions: (1). The conformal invariance of QCD in UV limit requires that $h(0) = f(0) = 1$. (2). The existence of a horizon with a nonzero Hawking temperature requires that $f(z) > 0$ for $0 < z < z_h$. With these restrictions, one can get the action for the quark pair and the inter-quark distance as

$$S(L, T) = \frac{\mathcal{T} R^2}{\pi \alpha'} \int_0^{z_0} dz \frac{h(z) \sqrt{f(z) + \dot{z}^2}}{z^2 \dot{z}}, \quad \dot{z} = \frac{dz}{dx} = \sqrt{\left( \frac{h(z) z_0^2}{h(z_0) z^2} \right) \frac{f^2(z)}{f(z_0)} - f(z)},$$

(38)

and

$$L = 2 \int_0^{z_0} \frac{dz}{\sqrt{\left( \frac{h(z) z_0^2}{h(z_0) z^2} \right) \frac{f^2(z)}{f(z_0)} - f(z)}}.$$

(39)

Likewise, the heavy quark potential can be found as

$$V = \frac{R^2}{\pi \alpha'} \int_0^{z_0} dz \frac{h(z) \sqrt{f + \dot{z}^2}}{z^2 \dot{z}} - \frac{R^2}{\pi \alpha'} \int_0^{z_{\text{max}}} dz \frac{h(z) \sqrt{f + \dot{z}_1^2}}{z^2 \dot{z}_1}, \quad 0 \leq z_0 \leq z_{\text{max}},$$

(40)
where
\[
\dot{z}_1 = \frac{dz}{dx} = \sqrt{\frac{h(z)z_{\text{max}}^2}{h(z_{\text{max}})z^2} \frac{f^2(z)}{f(z_{\text{max}})}} - f(z). \tag{41}
\]

One can see that the potential of Eq.(40) must go to zero at infinite separations (corresponding to \(z_0 = z_{\text{max}}\)). Also, one can check the results in some other holographic models such as the Sakai-Sugimoto model \[25\], one may find that the result is similar. Therefore, from this scheme the resulting potential seems to be negative for all quark-antiquark separations. Of course, if one doesn’t use this specific ansatz, the result may not show this behavior. Indeed, the resulting potential displays a kink-like screening if one uses the usual renormalization procedure \[36\].

IV. SUMMARY AND DISCUSSIONS

In this paper, we have studied the heavy quark potential in some deformed \(\text{AdS}_5\) models using a modified renormalization procedure. The quark-antiquark potential was obtained by calculating the Nambu-Goto action of string attaching the rectangular Wilson loop. It is found that the obtained potential is negative for all quark-antiquark separations.

Before discussing the result, let us first recall the usual renormalization procedure \[7\]
\[
V(L) = S(L, T) - 2S_0 \tag{42}
\]
where \(S_0\) is the self energy or the mass of each quark. We can see that the main difference between the two renormalization procedures Eq.(25) and Eq.(42) is the subtrahend in the numerator.

Why the heavy quark potential obtained from the two renormalization procedures is the same at zero temperature but differs at finite temperature or in some deformed \(\text{AdS}_5\) models? One possible reason is that at zero temperature, the total energy of the two quarks that separated by an infinite distance is consistent with the self energy of the two quarks due to the vacuum background. While in other backgrounds, the total energy of the \(\bar{Q}Q\) separated by an infinite distance may not only contain the self energy but also contain some energy from the field or the space such as the temperature, electric, etc.

Finally, we would like to discuss the reasonability of the results. The obtained potential in these deformed \(\text{AdS}_5\) models does not level off, implying the \(\bar{Q}Q\) pair never dissociates in a thermal medium. The results seem to be in contradiction with the experimental studies. However, in \[20\] the authors have analytically continued the string configurations into the complex plane. This manipulation yields a nonzero imaginary potential which would be responsible for melting the \(\bar{Q}Q\) \[37\]. So if we do this in the case of AdS/QCD, we may also get the imaginary potential. But due to the presence of the deformed warp factors in deformed \(\text{AdS}_5\) models, it is difficult to do this manipulation analytically and one has to resort to numerical methods. We hope to report our progress in this regard in future.

V. ACKNOWLEDGMENTS

We would like to thank Prof Hai-cang Ren for useful discussions. This research is partly supported by the Ministry of Science and Technology of China (MSTC) under the 973 Project no. 2015CB856904(4). Zi-qiang Zhang and Gang Chen are supported by the NSFC under Grant no. 11475149. De-fu Hou is supported by the NSFC under Grant no. 11375070, 11521064.

[1] J. Adams et al. [STAR Collaboration], Nucl. Phys. A 757, 102 (2005).
[2] K. Adcox et al. [PHENIX Collaboration], Nucl. Phys. A 757, 184 (2005).
[3] E. V. Shuryak, Nucl. Phys. A 750, 64 (2005).
[4] J. M. Maldacena, Adv. Theor. Math. Phys. 2, 231 (1998) [Int. J. Theor. Phys. 38, 1113 (1999)].
[5] S. S. Gubser, I. R. Klebanov and A. M. Polyakov, Phys. Lett. B428, 105 (1998).
[6] O. Aharony, S. S. Gubser, J. Maldacena, H. Ooguri and Y. Oz, Phys. Rept. 323, 183 (2000).
[7] J. M. Maldacena, Phys. Rev. Lett. 80, 4859 (1998).
[8] A. Brandhuber, N. Itzhaki, J. Sonnenschein and S. Yankielowicz, Phys. Lett. B 434, 36 (1998).
[9] S.-J. Rey, S. Theisen and J.-T. Yee, Nucl. Phys. B 527, 171 (1998).
[10] Y. Kinar, E. Schreiber and J. Sonnenschein, Nucl. Phys. B 566, 103 (2000).
[11] S.-x. Chu, D.f. Hou and H.-cang. Ren, JHEP 08, 004 (2009).
[12] Z.-q. Zhang, D. Hou, H.-c Ren and L. Yin, JHEP 1107 035 (2011).
[13] O. Andreev and V. I. Zakharov, JHEP 0704 100 (2007).
[14] S. He, M. Huang, Q.-s. Yan, Prog.Theor.Phys.Suppl.186 504 (2010).
[15] D. f. Zeng, Phys. Rev. D 78, 126006 (2008) [arXiv:0805.2733 [hep-th]].
[16] H. J. Pirner and B. Galow, Phys. Lett. B 679, 51 (2009) [arXiv:0903.2701 [hep-ph]].
[17] J. Greensite and P. Olesen, JHEP 9808, 009 (1998).
[18] F. Bigazzi, A. L. Cotrone, L. Martucci and L. A. Pando Zayas, Phys. Rev. D 71, 066002 (2005).
[19] L. Martucci, Fortsch. Phys. 53, 936 (2005).
[20] Javier L. Albacete, Yuri V. Kovchegov, and Anastasios Taliotis., Phys.Rev. D78 115007 (2008).
[21] O. Andreev and V. I. Zakharov, Phys.Lett.B, 645 437 (2007).
[22] S. He, M. Huang, Q.-s. Yan, Phys.Rev.D83 045034 (2011).
[23] J. Babington, J. Erdmenger, N. J. Evans, Z. Guralnik and I. Kirsch, Phys. Rev. D 69, 066007 (2004) [arXiv:hep-th/0306018].
[24] M. Kruczenski, D. Mateos, R. C. Myers and D. J. Winters, J HEP 0405 041.
[25] T. Sakai and S. Sugimoto, Prog. Theor. Phys. 113, 843 (2005); Prog. Theor. Phys. 114, 1083 (2006).
[26] S. He, M. Huang, Q. S. Yan and Y. Yang, Eur.Phys.J.C.(2010)66:187.
[27] O. Andreev, Phys. Rev. D 73, 107901 (2006).
[28] A. Karch, E. Katz, D. T. Son and M. A. Stephanov, Phys. Rev. D 74, 015005 (2006).
[29] J. P. Shock, F. Wu, Y. L. Wu and Z. F. Xie, JHEP 0703, 064 (2007).
[30] K. Ghoroku, N. Maru, M. Tachibana and M. Yahiros, Phys. Lett. B 633, 602 (2006).
[31] U. Gursoy and E. Kiritsis, JHEP 0802, 032 (2008).
[32] L. D. McLerran and B. Svetitsky, Phys. Rev. D24, 450 (1981).
[33] S. N. Nadiarni, Phys. Rev. D33, 3738 (1986).
[34] C. Gale and J. I. Kapusta, Finite-Temperature Field Theory (Cambridge University Press, 2006).
[35] J. Erlich, E. Katz, D. T. Son and M. A. Stephanov, Phys. Rev. Lett. 95, 261602 (2005).
[36] Y. Wu, D.f. Hou, H.-cang. Ren, [hep-ph/1401.3635].
[37] J. Noronha, A. Dumitru, Phys. Rev. Lett. 103 (2009) 152304.