Quarkonium in a non-ideal hot QCD Plasma

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Substantial anisotropies should occur in the hot expanding QCD plasma produced in relativistic heavy-ion collisions due to non-vanishing shear viscosity. We discuss the form of the real-time, hard thermal loop resummed propagator for static gluons in the presence of such anisotropies and the consequences for quarkonium binding. It has been predicted that the propagator develops an imaginary part due to Landau damping at high temperature. This should generate a much larger width of quarkonium states than the Appelquist-Politzer vacuum estimate corresponding to decay into three gluons. We argue that this might be observable in heavy-ion collisions as a suppression of the $\Upsilon(1S) \rightarrow e^+e^-$ process. Lastly, we consider the heavy quark (singlet) free energy just above the deconfinement temperature. In the "semi-QGP", $F_{Q\bar{Q}}(R)$ at distances beyond $1/T$ is expected to be suppressed by $1/N$ as compared to an ideal plasma.

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§1. The anisotropic QGP

This section deals with kinetic non-equilibrium effects in an expanding QCD plasma. During the first few fm/c when the temperature is high it is the longitudinal Bjorken expansion which matters most; for semi-peripheral collisions and/or close to the periphery of the fireball the transverse expansion may also be important but is neglected here for simplicity. As a consequence of the expansion, the particle momentum distribution in the local rest frame is anisotropic, if the scattering rate is finite: there is a net loss of particles with large $|p_z|$ from the fluid cell and redistribution of the momenta requires time. We assume that the momentum distribution can be parameterized as follows:

$$f(p) = f_{\text{iso}}(\sqrt{p^2 + \xi(p \cdot n)^2}) \simeq f_{\text{iso}}(p) \left[ 1 - \frac{\xi(p \cdot n)^2}{2pT} \left( 1 \pm f_{\text{iso}}(p) \right) \right].$$  (1.1)

The second expression is valid to first order in the anisotropy parameter $\xi$. $f_{\text{iso}}(p)$ is either a Bose or a Fermi distribution, respectively, and $n = e_z$ for longitudinal expansion. Such a correction $\delta f$ to the equilibrium distribution is also commonly employed in viscous hydrodynamics and has been argued to be consistent with azimuthal flow coefficients observed at RHIC (a transverse unit vector $n = e_T$ is used for calculating $p_T$ distributions). Near equilibrium one can relate moments of $\delta f$, which are proportional to $\xi$, to the shear.

Fig. 1 shows the time evolution of the ratio of the longitudinal to the transverse pressure in the central region of a high-energy heavy-ion collision (see, also, ref.3). This was obtained in ref.2 from a Boltzmann equation in relaxation time
approximation to all orders in the anisotropy parameter $\xi$. Near equilibrium, the assumed equilibration rate translates into a shear viscosity to entropy density ratio of $\eta/s = 0.1$. However, even for such “strong coupling” conditions large pressure anisotropies are caused by the rapid longitudinal expansion during the early, high-temperature stage of a heavy-ion collision ($\tau \lesssim 3$ fm/c). Note that $p_L < p_T$ does not suppress transverse hydrodynamic flow but may reflect in the system-size and energy dependence of the transverse energy $dE_T/dy$ in the final state.

1.1. Real part of the quarkonium potential

Anisotropies of the momentum distributions affect various processes in the plasma. In particular, rotational symmetry in the local rest frame is broken and the screening length acquires an angular dependence. At high temperature and in the weak-coupling approximation it can be obtained explicitly from the “hard thermal loop” (HTL) resummed propagator for static electric gluons. To linear order in $\xi$.

$$\text{Re} \Delta^{00}(p) = \frac{1}{p^2 + m_D^2} \left( 1 - \xi m_D^2 \frac{2}{p^2 + m_D^2} - \xi m_D^2 \frac{(p \cdot n)^2 / p^2}{p^2 + m_D^2} \right).$$  \hspace{1cm} (1.2)

Here, $m_D = gT\sqrt{N_c}/3$ denotes the standard Debye mass of the equilibrium plasma. The one-gluon exchange potential follows essentially from the Fourier transform,

$$\text{Re} \ V(r) = V_{\text{iso}}(r) \left( 1 + \xi \left[ \frac{\hat{r}^2}{6} + \frac{\hat{r}^2}{48} + \frac{\hat{r}^2}{16} \cos(2\theta) \right] \right).$$  \hspace{1cm} (1.3)

Here, $\hat{r} \equiv r m_D$, $\cos \theta \equiv \hat{r} \cdot \hat{n}/r$,$r$, and $V_{\text{iso}}(r) = -\frac{g_s C_F}{r} \exp(-\hat{r})$ is the well-known Debye-screened Coulomb potential. The potential \hspace{1cm} \hspace{1cm} \hspace{1cm} (1.3) reduces thermal screening effects as compared to an ideal plasma in local equilibrium ($\xi = 0$). On the other hand, one can absorb this $\xi$ dependence to a large extent by a redefinition of the hard scale $T(\xi)$ and/or of the Debye mass $m_D(\xi)$. The most useful approach with respect to applications to heavy-ion collisions would probably amount to adjusting the initial $T_0$
in such a way that the entropy (per unit of rapidity) in the final state remains fixed as $\xi_0$ is varied. The resulting implicit dependence $T_0(\xi_0)$ can be determined only from solutions of viscous hydrodynamics or transport theory (with gluon-number changing processes).

For semi-realistic estimates of the binding energies of charmonium and bottomonium states one needs to add the linear confining potential $\sim \sigma r$, where $\sigma \simeq 1$ GeV/fm is the SU(3) string tension. Its temperature dependence to be used in the real-time formalism must presently be modelled. In the phenomenologically relevant range $T/T_c = 1 - 3$ the “interaction measure” $\langle e - 3p\rangle/T^4$ in QCD is large. The potential at infinite separation is thus sometimes modelled as $V_\infty(T) \simeq 2a/T$ with $a \approx 0.08$ GeV$^2$ a constant of dimension two. Such a model has in fact been proposed long ago in ref.\cite{13}

$$V(r) = \left[ -\frac{\alpha_s C_F}{r} + 2\frac{\sigma}{m_D}(e^r - 1) - \sigma r \right] e^{-r}, \quad (1.4)$$

which also provides a smooth interpolation to short distances. In this model the temperature dependence of the binding energy $E_{\text{bind}} = \langle \Psi | \hat{H} - V_\infty | \Psi \rangle - 2m_Q$ of small bound states such as the $\Upsilon$ actually turns out to originate mostly from $V_\infty(T)$\cite{12}.

This is confirmed by the spectral function of pseudo-scalar bottomonium published in ref.\cite{10} which has been replotted on a logarithmic scale in fig. 2: while the ground state peak shows little temperature dependence, the continuum threshold decreases rapidly as $T$ increases. Similar spectral functions have been shown in ref.\cite{13}.

The anisotropy or viscosity affects the excited states more strongly than the compact 1S bottomonium state. In this model the binding energy at $T/T_c = 1.1$, for example, increases by about 50%. Alternatively, if a “dissolution” temperature is defined from $|E_b| = T_{\text{dis}}$ then this increases from $\simeq 1.15T_c$ when $\xi = 0$ to about $1.35T_c$ when $\xi = 1$. Thus, excited states should melt less easily if the QGP exhibits
an anisotropy during the early stages of the expansion. Lastly, there is a splitting of the states with \( \mathbf{L} \cdot \mathbf{n} = 0 \) and \( \mathbf{L} \cdot \mathbf{n} = \pm 1 \), respectively, since rotational symmetry is broken. Understanding the properties of excited states of the \( \Upsilon \) is important for phenomenology because they contribute through feed down to the yield of the 1S state \(^{13}\).

1.2. Imaginary part of the quarkonium potential

At finite temperature, the quarkonium potential also acquires an imaginary part \(^{13}\) at order \( g^2 C_F \) due to Landau damping of the exchanged nearly static gluon. It can again be obtained from the Fourier transform of the HTL-resummed real-time propagator (the “physical” component of the Schwinger-Keldysh representation) for static \( A_0 \) fields. Taking the imaginary part corresponds to cutting open one of the hard thermal loops of the HTL propagator and can be viewed, microscopically, as the scattering of the space-like exchanged gluon off a thermal gluon \(^{15}\) \( g + (Q\bar{Q}) \rightarrow g + Q + \bar{Q} \). To order \( \xi \) and in the leading log \( 1/\hat{r} \) approximation \( \text{Im} V \) is given by \(^{10}\)

\[
\text{Im} V(r) = -\frac{g^2 C_F T}{4\pi} \hat{r}^2 \log \frac{1}{\hat{r}} \left( \frac{1}{3} - \frac{3 - \cos(2\theta)}{20} \right). \tag{1.5}
\]

This corresponds to a decay width for a Coulomb ground state of

\[
\Gamma = \frac{16\pi T m_Q^2}{g^2 C_F M_Q^2} \left( 1 - \frac{\xi}{2} \right) \log \frac{g^2 C_F M_Q}{4\pi m_D}. \tag{1.6}
\]

For \( \xi = 0 \) one can in fact evaluate the matrix element of \( \text{Im} V \) between Coulomb wave functions without resorting to the log \( 1/\hat{r} \gg 1 \) approximation,

\[
\Gamma = \frac{T}{\alpha_s C_F M_Q^2} \frac{m_Q^2}{(1 - \kappa^2)^2} + 4 \log \frac{1}{\kappa} \kappa = \frac{1}{\alpha_s C_F M_Q}. \tag{1.7}
\]

The width obtained from \(^{17}\) is shown in fig. 3. For temperatures accessible to the RHIC and LHC colliders, \( \Gamma_\Upsilon \) is on the order of \( 20 \text{ MeV} - 100 \text{ MeV} \). This
can be compared to the $\Upsilon \to e^+e^-$ decay width in vacuum which arises from the annihilation of the $b, \bar{b}$ quarks into di-electrons $^7$: $\Gamma_{\Upsilon \to e^+e^-} = 1.34 \text{ keV}$. Because the electromagnetic decay width is much smaller than the inverse lifetime of the QGP the actual $\Upsilon$ peak observed in the di-lepton invariant mass distribution would not be broadened as compared to vacuum.

Nevertheless, $\Upsilon$ states which have been broken up into individual $b, \bar{b}$ quarks reduce the yield of di-leptons in the peak. One of the contributions to this process is $g + (Q\bar{Q}) \to Q + \bar{Q}$ dissociation $^{19}$ by a thermal gluon. However, it has been found that this process leads to rather small $\Upsilon$ dissociation rates. Indeed, ref. $20$ argues that in the limit of loose binding, $|E_b| \ll T$ due to strong screening of the attractive interaction, it becomes more efficient to scatter a thermal gluon off one of the quasi-free $b$-quarks thereby breaking up the bound state. The width obtained from the imaginary part of the HTL propagator indicates that damping of the exchanged gluon in the heat bath also provides a large contribution to the thermal $\Upsilon \to b + \bar{b}$ rate.

At low temperature, once pions have formed, a non-zero (but exponentially small) width emerges due to $\pi + \Upsilon \to B + \bar{B}$ corresponding to tunneling of the $b, \bar{b}$ quarks in the background field of the pion $^{21}$.

It is certainly interesting to compare to a strongly coupled theory. Using the gauge-gravity duality, the static potential (or Wilson loop) $^{22}$ and thermal effects at short distances $^{23}$ have been computed in $N = 4$ supersymmetric Yang-Mills at large (but finite $^{24}$) t’ Hooft coupling $\lambda = g^2N$ and $N \to \infty$. At $T = 0$,

$$V_{Q\bar{Q}}(r) = -\frac{4\pi^2}{\Gamma(1/4)^4} \frac{\sqrt{\lambda}}{r}. \quad (1.8)$$

The $\sim 1/r$ behavior follows from conformal invariance of the theory. Also, the potential is non-analytic in $\lambda$. Clearly, the coupling should not be very large or else the properties of the resulting bound states are qualitatively different from the $\Upsilon$ etc. states of QCD (numerically, $4\pi^2/\Gamma(1/4)^4 \approx 0.23$).

At finite temperature, the potential develops an imaginary part when the string dangling in the fifth dimension (with end points on our brane) approaches the black hole horizon which sets the temperature scale and “melts” $^{23}$. Fluctuations near the “bottom” of the string should generate such an imaginary part of the Nambu-Goto action at even lower temperatures already $^{25}$

$$\Gamma_{Q\bar{Q}} = -(\psi|\text{Im } V_{Q\bar{Q}}|\psi) \simeq \frac{\pi\sqrt{\lambda}}{48\sqrt{2}a_0} \left[ 45 \left( a_0 T b \right)^4 - 2 \right], \quad (1.9)$$

where $|\psi\rangle$ denotes the unperturbed Coulomb ground state wave function, $a_0 = \Gamma(1/4)^4/2\pi^2\sqrt{\lambda} m_Q$ is the “Bohr radius” for the Maldacena potential $^{18}$ and $b = 2\Gamma(3/4)/\sqrt{\pi} \Gamma(1/4) \approx 0.38$ is a numerical constant. Here the width decreases with the quark mass and with the t’ Hooft coupling approximately as $\Gamma_{Q\bar{Q}} \sim 1/\lambda m_Q^3$.

$^1$ It is therefore proportional to the square of the quarkonium wave function at the origin, similar to the hadronic decay into three gluons discussed by Appelquist and Politzer$^{17}$ but unlike the thermal QCD width discussed above.
it increases rapidly with the temperature, \( \sim T^4 \). For \( m_Q = 4.7 \text{ GeV}, \ T = 0.3 \text{ GeV}, \ \sqrt{A} = 3 \) one finds \( \Gamma_Y \simeq 50 \text{ MeV} \). The \((\bar{Q}Q) \rightarrow (\bar{q}Q)(\bar{q}q)\) breakup due to string splitting at has been considered in ref.\(^{27}\)

The suppression of the yield of di-leptons from \( Y(1S) \) decays in the final state should be significant.\(^{20}, 26\) Neglecting “regeneration” of bound states from \( \bar{b} \) and \( b \) quarks in the medium, the number of \( Y \) mesons in the plasma which have not decayed into unbound \( b \) and \( \bar{b} \) quarks up to time \( \tau \) after the collision is

\[
N(t) \simeq N_0 \exp \left( -\int_{\tau_0}^{\tau} dt \Gamma_Y(t) \right) .
\]

This solution assumes that \( \Gamma_Y(T(t)) \) is a slowly varying function of time. The initial number of \( Y \) states may be estimated from the multiplicity in \( p+p \) collisions times the number of binary collisions at a given impact parameter: \( N_0 \simeq N_{\text{coll}} N_{\text{pp}} \). Thus, the “nuclear modification factor” \( R_{AA} \) for the process \( Y \rightarrow \ell^+\ell^- \) is approximately given by \( R_{AA}(Y \rightarrow \ell^+\ell^-) \simeq \exp(-\bar{\Gamma} \tau) \), where \( \bar{\Gamma} \) denotes a suitable average of \( \Gamma(T) \) over the lifetime of the quark-gluon plasma. Due to the strong temperature dependence of the width, this average is dominated by the early stage. Experimental measurements of \( R_{AA}(T \rightarrow \ell^+\ell^-) \) at RHIC appear to indicate a suppression but so far it has not been possible to disentangle various \( Y \) states.

\( \S 2. \) The static gluon propagator in the “semi-QGP”

This section is about the HTL propagator in Euclidean time for temperatures just above deconfinement. Hidaka and Pisarski suggested that for \( T \) not far above \( T_c \) hard modes in the QCD plasma may still be weakly coupled but propagating in a non-perturbative background \( A_0 \) field.\(^{29}\) The expectation value of the Polyakov loop is parameterized as

\[
L = \exp \left( -\frac{Q}{T} \right) , \quad Q \equiv g A_0 , \quad \ell = \frac{1}{N} \text{tr} \ L .
\]

At high \( T \) the eigenvalues of the matrix \( L \) approach 0 and so its normalized trace

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig4}
\caption{Left: Illustration of “repulsion” of the eigenvalues of the Polyakov loop for \( N = 3 \) colors. Right: Singlet free energy versus distance \( R \), data from.\(^{32}\)}
\end{figure}
$\ell \to 1$. On the other hand, as $T \to T_c$, the eigenvalues “repell” and eventually, at \( T_c - 0 \), distribute uniformly over a circle such that $\ell = 0$. Note that $\ell(T_c + 0) \approx 0.4$ in SU(3) Yang-Mills, far from unity, and hence that $A_0 \sim T/g$ is non-perturbatively large. It is in this sense that the “semi-QGP” just above $T_c$ is a weakly coupled but non-perturbative phase.

Gluons which couple to the background field acquire an additional “mass”; the static propagator in the semi-QGP (summed over colors) becomes, schematically,

$$\frac{1}{N} \langle \text{tr} \mathbf{L}(R) \mathbf{L}(0) \rangle \sim g^2 C_F \int \frac{d^3 p}{(2\pi)^3} e^{ipx} \sum_{kl} \frac{1}{p^2 + m_D^2 + (Q_l - Q_k)^2}. \quad (2.2)$$

The $N$ diagonal gluons are screened only over large distances of order $1/m_D$, where $m_D^2 = O(g^2 T^2 N)$ is the usual Debye mass. The $N^2 - N$ off-diagonal gluons corresponding to $l \neq k$, however, acquire a mass of order $T$ when $A_0^2 - A_0 \sim T/g$. The heavy-quark free energy $F_{QQ}(R)/T$ defined from the correlator of Polyakov loops (which may not necessarily be identified with the potential in real time) is expected to show the following qualitative behavior (assuming, for simplicity, large $N$):

$$RT \ll 1 : \quad F_{QQ}(R) \sim -\frac{g^2 N^2}{NR} \sim \frac{g^2 N}{R} \quad (2.3)$$

$$1 \ll RT \ll \frac{1}{g} : \quad F_{QQ}(R) \sim -\frac{g^2 N}{NR} \sim \frac{g^2}{R} \quad (2.4)$$

$$\frac{1}{g} \ll RT : \quad F_{QQ}(R) \sim -\frac{g^2 N}{NR} e^{-m_D R} \sim \frac{g^2}{R} e^{-m_D R}. \quad (2.5)$$

In the intermediate region \( (2.4) \) the $N^2 - N$ heavy off-diagonal gluons have decoupled and $F_{QQ}$ is suppressed by $1/N$ as compared to its behavior at short distances. In fig. 4 we show the singlet free energy for SU(3) Yang-Mills just above $T_c$ obtained numerically on $32^3 \times 4$ lattices. Indeed, it does appear to show a two-slope behavior with a transition at $RT \approx 2$. On the other hand, near $T_c$ the large distance behavior could be affected by finite volume artifacts simulations on larger lattices may hopefully become available in the future. If the behavior seen in the present data is confirmed then this may indicate that just above $T_c$ the properties of large, excited quarkonium states may be modified as compared to predictions from a standard potential model such as \( (1.4) \).

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