Signatures for strongly coupled Quark-Gluon Plasma

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Abstract. Dramatic changes had occurred with our understanding of Quark-Gluon Plasma, which is now believed to be rather strongly coupled, sQGP for short. Hydrodynamical behavior is seen experimentally, even for rather small systems (rather peripheral collisions). From elliptic flow the interest is shifting to even more sophisticated observable, the conical flow, created by quenched jets. The exact structure of sQGP remains unknown, at the moment the best picture seem to be a liquid made partly of binary bound states. As we discuss at the end, those can be possibly seen in the dilepton spectra, as “new vector mesons” above $T_c$.

1. Introduction: a strongly coupled QGP

Although some speculations about further running of the QCD coupling constant at $T > T_c$ to higher values were around for a while (see e.g. Fig.5 in [1]), the main arguments for a new QGP regime at $T = (1 - 3)T_c$ are: (i) a triumph of hydrodynamical description of RHIC collisions, explaining strong radial and elliptic flows [2, 3]; (ii) discovery of multiple bound states of heavy and light quarks at $T > T_c$, from the lattice correlators [4] and directly from Schreodinger/Klein-Gordon/Dirac equations [5, 6] with properly defined (lattice-based) potentials [7]. These general arguments are by now well accepted in the community, see e.g. [8] and the “experimental white papers”, and so I would not discuss them here in any detail, just make a couple of introductory remarks.

One is a reply to the question raised in the opening talk by J.P.Blaizot, who wandered why do we know that equilibration happens as early as at time $\tau \sim 1/2 fm$. The most direct answer to that is that hydro works well till very peripheral collisions, abruptly failing only when the initial “almond” is less than 1 fm wide.

Another comment is on an inevitable question: How strong is strong enough? The answer, suggested in a “strong coupling” paper [5], is that it should be sufficiently strong to bind quasiparticles together. If so, lines on the phase diagram were bound states gets marginal (have zero binding) separate the sQGP region. Presumably, the most stable is the most attractive $gg$ singlet channel. These lines are not singularities of the thermodynamical quantities: as marginal states become virtual the scattering amplitude remains continuous.

One particular bound state, $J/\psi$, got special attention. According to lattice results [4] using MEM method and our calculations based on lattice potentials [6] it does not melt for $T < 2.5 - 3T_c$.

An important new feature sQGP scenario [6] is and possible existence of hundreds of colored bound states, such as $gg_3$ and $(gg)_8$, which outnumber the colorless ones by about 10:1. This
prediction is not yet tested on the lattice, but it will be. Calculations presented in that work solved the “pressure puzzle”, explaining where additional contribution to that of (not-so-light) quasiparticles comes from. New binary bound states help to get the EoS right, see [6].

This is of course only a beginning of the road toward a full understanding of sQGP. Studies related to 3-body physics (known as Efimov effect) have revealed that if there are marginal 2-body states, there must also be 3-body ones, some of them even deeply bound. A really manybody studies are not yet done, and it may well be that sQGP is polymerized into some chains, or crystallize into a (locally ordered) liquid. Color degrees of freedom are numerous and provide lots of possibilities, so a lot of hard work is ahead to understand the sQGP structure.

As experience of other fields shows, global EoS is probably too insensitive a measure to tell the difference between all those models, at least with the level of accuracy we now have. Transport properties like viscosity is a different matter, they are known to be quite sensitive to coupling and can vary by several orders of magnitude for known liquid and plasmas. Its theory is still missing, and a lot of efforts would be needed to have it.

One early idea suggested in [5] is based on the marginal states with small binding, which may be related to large scattering length and hydro behavior. This idea is further confirmed by experiments with cold atoms, where precisely this mechanism – known as the Feshbach resonance – is used. Not only the elliptic flow in agreement with hydrodynamics was observed [9], but more recent experiments such as [10] have revealed that in fact frequency of the two lowest (z and r-modes) of oscillations in elliptic trap both agree with hydro prediction better than a percent. Furthermore, the damping-to-frequency ratio diverges from ~ 1 to about 10^{-3} near the Feshbach resonance. It means that there hydro oscillations may be repeated about a thousand times, before dissipative viscose effects take over and kill it. And all of it happens, let me repeat, where the interaction is the strongest, corresponding to viscosity as strikingly small [11], as it is for QGP at RHIC.

Another connection on which I would like to provide an update is a connection to $\mathcal{N}=4$ SUSY YM theory and AdS/CFT correspondence. Let me just remind that this theory has zero beta function and the coupling does not run, and for all values from 0 to infinity the finite-T version of it is in the same Coulomb (QGP-like) phase. In [12] we argued that matter is made not from quasiparticles, which have large masses $M_{q,g} \sim \sqrt{\lambda T}$ (where $\lambda = g^2 N_c >> 1$ is ’t Hooft gauge coupling) but of deeply bound binary states of those with much smaller masses $M_{\text{mesons}} \sim T$ which we found for relativistically rotating states with large angular momentum. Recently there was a progress in introducing fundamental massive quarks into AdS/CFT theory, living on a separate D7 brane, with results for “charmonium” spectroscopy in strong coupling [13]. It was found that the “s-wave” states with zero orbital momentum do not “fall on the center” but survive and have masses $M_{\text{mesons}} \sim M_q/\lambda^{1/2}$, the same scale as needed to populate the heatbath (as advocated in our paper). Thus a splitting between $J/\psi, \eta_c$ and $\chi$ states is large and parametric in strong coupling. Some dreams about using the AdS/CFT correspondence for dynamical description of high energy collisions start to be discussed.

2. Where does the energy of the quenched jets go?
Quenching of high-$p_t$ QCD jets leads to deposition of large energy into matter. If it is a near-perfect non-dissipative liquid, this energy should propagate outward as a “conical flow” similar to well known sonic booms from supersonic planes. This idea is developed in a recent paper [14].

In what follows, we would like to treat separately two different types of energy losses: (i) the radiative losses, producing mostly relativistic gluons and (ii) the scattering/ionization losses,

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1. It is based on expected Casimir scaling of forces in all channels.
2. As I learned in Kolkata, it was also suggested earlier by H.Stocker, see his talk for details.
Figure 1. (a) A schematic picture of flow created by a jet going through the fireball. The trigger jet is going to the right from the origination point (the black circle at point B) from which sound waves start propagating as spherical waves (the dashed circle). The companion quenched jet is moving to the left, heating the matter and thus creating a cylinder of additional matter (shaded area). The head of the jet is a “non-hydrodynamical core” of the QCD gluonic shower, formed by the original hard parton (black dot A). The solid arrow shows a direction of flow normal to shock cone at the angle $\theta_M$, the dashed arrows show the direction of the flow after shocks hit the edge of the fireball. (b) Distribution of the radial flow velocity $v_r$ in the $x - r$ plane (units used are fm) for a jet propagating through the diameter of Au. We used $dE/dx = 1 \text{ GeV/fm}$ and the width of the Gaussian $\sigma = 1 \text{ fm}$.

Figure 1 explains a view of the process, in a plane transverse to the beam. Two oppositely moving jets originate from the hard collision point B. Due to strong quenching, the survival of the trigger jet biases it to be produced close to the surface and to move outward. This in turn forces its companion to move inward through matter and to be maximally quenched.

The energy deposition starts at point B, thus a spherical sound wave appears (the dashed circle in Fig.1(a) ). Further energy deposition is along the jet line, and is propagating with a speed of light, till the leading parton is found at point A at the moment of the snapshot. As is well known, the interference of perturbations from a supersonically moving body (such as a supersonic jet plane or meteorite) create conical flow behind the shock waves. The angle $\theta_M$ defined in the figure is given by a simple geometrical condition: distance $AB = ct$ while $CB$ is the distance travelled by a shock. Since the velocity of the shock depends on its intensity,
the cone should in fact be somewhat rounded near its top: that is ignored in the figure. The region near the head of the jet, which we would refer to as a “non-hydrodynamical core”, which is constantly producing gluons, which emit new ones etc: the whole shower is a complicated nonlinear phenomenon which should obviously be treated via tools of quantum field theory. The energy/momentum flow through the core boundary we will treat phenomenologically, identifying it with $dE/dx$ of the second type, which we will approximate by a time-independent constant. This should thus lead to stationary-state conical flow, depending on $x - t$ only.

**Hydrodynamical equations** we use are linear since the whole energy deposited by the jets can be treated as small compared to the total energy of the medium. We will use cylindrical coordinates with $x$ along the jet axis and $r$ orthogonal to it. For simplicity we will assume that the perturbed medium is homogeneous and at rest. In this linearized approximation the nonzero components of the stress-energy tensor are

$$
\delta T^{00} = \delta \epsilon \quad \delta T^{0i} = (\epsilon + p)v^i \quad \delta T^{ij} = \delta p \delta^{ij}
$$

where $\epsilon$ and $p$ are the internal energy and pressure and $\vec{v}$ is the velocity field of the perturbation. Thus, by recalling that $\frac{\partial p}{\partial \epsilon} = c_s^2$ (with $c_s$ velocity of sound) the energy and momentum conservation equation $\partial_\mu \delta T^{\mu \nu}$ can be written as

$$
\begin{align*}
\partial_0 \delta T^{00} + \partial_i \delta T^{0i} &= 0 \\
\partial_0 \delta T^{0i} + c_s^2 \partial_i \delta T^{00} &= 0
\end{align*}
$$

The initial conditions are set by the the process of thermalization of the energy lost by the jet. As already mentioned, this thermalization process is very complicated and should take place at distances of the order $\Gamma_s$ from the production point. As $\Gamma_s$ is also the typical size of the liquid cells, we will simply consider that there is a variation of the energy and momentum at the position of the particle.

This initial condition can be easily expressed if we first concentrate in the perturbation due to the propagation of the particle in an interval $dt$ around its position. In this case the previous conditions lead to

$$
\begin{align*}
\delta T^{00}(t_0, x) &= \Delta E \delta^3(\vec{x}(t_0) - \vec{x}) \\
\delta T^{0x}(t_0, x) &= \Delta P \delta^3(\vec{x}(t_0) - \vec{x})
\end{align*}
$$

where we have assumed that the particle moves in the $x$ direction at the speed of light ($c = 1$) and $\vec{x}(t)$ is the trajectory of the particle ($x = t$). We assume also that the energy and momentum loss are equal and, as we look for a constant drag, we set $\Delta E = \frac{dE}{dx} dt$. Assuming that $c_s$ is constant, it can be shown that the solution for these conditions is given by:

$$
\delta T^{00} = \frac{1}{c_s} \frac{dE}{dx} dt_0 \left( \partial_t - \partial_x \right) \Theta(t - t_0)
$$

$$
\frac{[\delta (R - c_s(t - t_0)) - \delta (R + c_s(t - t_0))]}{4\pi R}
$$

Where $R = |\vec{x} - \vec{x}(t_0)|$, the distance from the observation point to the emission point at the intersection of the past “sound cone” and a jet world line. The reader can readily identify the argument of the derivatives as the retarded Green function of the wave equation, and be warned that $R$ depends on observation time and space point. The calculation is close to that done for electromagnetic waves (Cerenkov radiation), except that sound is a longitudinal excitation.
In general, let a jet be born at the initial point $t_i$ and dies the final point $t_f$. A point-like source creates two spheres of perturbation, with the Mach cone tangent to both:

$$\delta T^{00} = \frac{1}{c_s^2} \frac{dE}{dx} \left( \frac{\delta(t - t_i - R(t)/c_s)}{4\pi R(t_i)} - \frac{\delta(t - t_f - R(t)/c_s)}{4\pi R(t_f)} \right)$$

$$- \frac{2}{c_s^2} \frac{dE}{dx} \frac{\partial}{\partial x} \Theta(t - t_i - R(t)/c_s) \Theta(t_f - t + R(t)/c_s)$$

when integrating over the whole space, the two spherical terms give the total deposited energy $\frac{dE}{dx}(t_f - t_i)$. The conical term apparently has no energy. In order to explain why, we prefer to regularize the problem, introducing a distributed source, e.g. a Gaussian one, see the paper for details.

**How can the effect be observed?** The most important feature of conical flow is its direction, at the so called Mach angle (6). Let us start with the simplest case of central collisions, in which the cross section of the fireball is a circle and the elliptic flow is absent. Most quenched jets going through it will pass a length comparable to its diameter, and thus the high-energy jets which are not quenched completely and reach through the fireball will produce conical flow during time $\tau \sim 10 - 15 \text{fm/c}$. The angle $\theta_M$ defined in the figure is given by a simple geometrical condition: the distance travelled by jet during the proper time interval.

$$\cos \theta_M = \frac{1}{c_s} \int c_s dt$$

The appropriate value for the speed of sound is a weighted average of that for QGP ($c_s \approx 1/\sqrt{3}$), the mixed phase $c_s \approx 0$ and a “resonance gas”, $c_s \approx \sqrt{2}$, $\bar{c}_s \approx .33$. We thus conclude that in this case the emission angle of the conical flow should be at the angles about 70 degrees relative to the jet, or in radians

$$\theta_{\text{emission}} = \pi \pm \arccos(\bar{c}_s/c) \approx 1.9, 4.3$$

In Fig.2(left) we show preliminary data from STAR experiment. The peak near angle zero consists of particles from a triggered outward moving jet, while particles from the companion jet should result in a peak near $\Delta \phi = \pi$. Such peak is clearly seen in pp collisions (a), in which no matter is present: but in AuAu collision (b) one finds instead a double-peaked distribution with a minimum at the original jet direction $\Delta \phi = \pi$. Remarkably, as Fig.2(b) shows, this value $\pi \pm 1.2$ nicely agrees with the observed positions of the two maxima of the distribution of the secondaries, relative to the direction of the triggered hard charged hadron.

At this conference B.Jacak have for the first time reported analogous data from PHENIX experiment (see her talk). It selects companions with $p_t > 1 \text{GeV}$ and different centrality classes. Five out of six of them (except the most peritheral bin) show even sharper peaks at the same angle as we predict. Furthermore, after the conference STAR reported spectra of associated partiles, measuring in particular their mean $<p_t>($ $\phi$). It also has maxima in the direction away from $\phi = \pi$, where one finds a clear a minimum, consistent in magnitude with the mean transverse momentum in average background events. Thus, all known facts strongly support our explanation³ of the peaks.

³ It has been argued at this conference by I.Vitev (see his talk) that radiation effects calculated with Landau-Pomeranchuck-Migdal effect can also produce a maximum at large angles, similar to Mach angle we discuss. Those calculations however used static (infinitely massive) scatterers. If this unrealisitic approximation is lifted, a maximum presumably corresponds to an angle 90 degrees in CM. To make it as large as 70 degrees observed, one would need a combined mass of all target partons participating in the collision to be $M \sim 10 \text{GeV}$ or so, which is hard to imagine.
Figure 2. (left) The background subtracted distribution of secondaries in azimuthal angle $\Delta \phi$ which are associated with triggered hard charged hadron $p_t > 4 \text{GeV}$, by STAR collaboration [20]. The figure (a) is for $p+p$ and (b) for 5% most central Au+Au collisions. (right) Schematic $T$-dependence of the masses of $\bar{q}q$ states. $A, V, S$ and $PS$ stand for axial, vector, scalar and pseudoscalar states. The dash-dotted line shows a behavior of twice the quasiparticle mass. Two black dots indicate places where we hope the dilepton signal may be observable.

3. Direct observation of binary bound states in sQGP via dileptons?

Photons and dileptons are “penetrating probes”, providing information about all stages of heavy ion collisions [21]. It was then assumed that the invariant mass $(M^2 = (p_{l+} + p_{l-})^2)$ spectra of penetrating probes produced in QGP should be monotonously decreasing with the dilepton mass, contrasting with peaks from hadronic stage due to familiar vector mesons $\rho, \omega, \phi, J/\psi$,...

Now the logic of it is changed [19]: if meson-like bound states continue to exist at $T > T_c$, one may see them in dilepton spectra as well. Furthermore, even at high $T$ when there are no bound states, there still exists a near-threshold enhancement, which may still be detected and used to infer the value of the quasiparticle masses and the strength of their interaction.

The observability of new states is obscured by the fact that their mass is “sliting” with $T$, from very large masses $M \sim 1.5 - 2 \text{GeV}$ where the states gets unbound to normal values at freezeout. One may thus wander if the very existence of the peaks can be observed. We hope it can be done in two endpoints mass regions:

(i) $M \sim 1.5 - 2 \text{GeV}$ corresponding to $T \approx T_c(\text{zero binding})$ and the edge of the quasiparticle continuum, at the initial QGP at RHIC.

(ii) $M \sim 0.5 \text{GeV}$, where we expect to see the contributions of the modified vector mesons at $T \approx T_c$.

The basic idea is very simple: the probability of $\bar{q}q \rightarrow l^+l^-$ process is enhanced by the initial state attractive interaction. Attractive interaction obviously correlate $\bar{q}$ and $q$ in space and increases the chances to find $\bar{q}q$ close to each other and annihilate. In such general form, the enhancement persists whether the potential is deep enough to make bound levels or not, and whether quark quasiparticles have small or large width.

In the case of pure Coulomb interaction the relevant parameter $z = \pi \alpha / v$ contains the ratio of the coupling strength to velocity. The enhancement is given by the well known Gamow factor$^4$

$^4$ The sign in exponent is the opposite to that on the original Gamow factor, for alpha-particle interaction with
Figure 3. Modification of the spectral density versus the invariant mass in $M_q$ units for different temperatures: (a) 1.2 $T_c$ (cross) and 1.4 $T_c$ (circle), (b) 1.7 $T_c$ (cross) and 3 $T_c$ (circle) and the correction due to quark mass (line).

\[ F_{\text{Gamow}} = \frac{z}{1 - \exp(-z)}. \] Note that the result is $\sim 1/v$ at small $v$ (large $z$) and 1 at small $z$. It cancels the velocity in the phase space and makes the production rate to jump to a finite value at the threshold.

New (sQGP) part of this picture is at $T > T_c$, which starts and ends at points marked by small black circles. Two of the black dots, at the mixed phase $T = T_c$, have a benefit of long time spent there during expansion. The same is true for the third black point, due to near-constancy of the mass of weakly bound states at $T = 1.5 - 2T_c$. The near-threshold bump is at about the same location at higher $T$ as well. Thus one may hope that the structures corresponding to these endpoints (black circles at Fig.2(right) may be detected.

In it we have shown only states made of $u, d$ quarks, ignoring the strange one. Needless to say, those exist and can be also produced. The peaks of $\bar{s}s$ $\phi$-like states should be shifted in mass upward by $2m_s \approx 250 - 300$ MeV, but their contribution to dilepton spectra is proportional to square of the electric charge $q_s^2 = 1/9$, which is 3 times smaller than the average of $q_u^2 = 4/9, q_d^2 = 1/9$. Strange states are more promising for pseudoscalar/scalar signals to which we turn below.

Omitting technical part, let me just say that non-relativistic Green’s function $G_{E+\Gamma_i}(r, \bar{r})$ obeys the usual Schrodinger equation:

\[
[-\frac{1}{m} \nabla^2 + V(r) - (E + i\Gamma)]G_{E+\Gamma_i}(r, \bar{r}) = \delta^3(r - \bar{r})
\]

with inter-particle potential at relevant $T V(r)$ extracted from the lattice. The benefit of using this method is that one avoid summations over levels, and bound states are automatically included together with scattering ones.

Let us now explain our units. Since the effective mass of quark quasiparticle in the temperature interval considered is not known accurately, we use it as our basic energy unit. In plots twice this value appear as a threshold, to which we tentatively ascribe to it a particularly simple value $2M_q = 2$ GeV: the reader should however be warned that this is a “GeV” in nuclei or that used for HBT correlations, corresponded to a repulsive Coulomb force.
quotation marks, to be rescaled appropriately later when the value of quark effective masses in QGP will be better known.

What we learned from these calculations is how the spectral density changes as the bound states become less and less bound until the system arrives to zero binding (zero binding point). In figure 3 we show the shape of the dilepton spectral density for different temperature (1.2 $T_c$, 1.4 $T_c$, 1.7 $T_c$, 3 $T_c$). One can observe how exactly enhancement in the bound state and threshold region changes. The height of this enhancement depends on the proximity of the level; at 1.7 $T_c$, where the bound level is close to threshold, we observe a big modification of the spectral function at 2 GeV. Note that proximity of the level to threshold happens in rather wide $T$ interval, roughly between 1.5 $T_c$ to 2 $T_c$ (the zero binding point [5]. Note also that the enhancement is still seen at temperatures as high as 3 $T_c$, when all bound states have already been melted.

The issue of the observability of these peaks depend on their width: its magnitude is being addressed right now and I am not yet ready to report those. The pictures above only have “numerical width” for plotting, and thus are not realistic. However it is very likely that the width is not overwhelming to preclude their observation.

Acknowledgments.

I am thankful to organizers for making this important conference such a success. Wish to express once again my gratitude to Bikash Sinha for all his efforts over the years, and wish him to go on with the same energy for many years to come. I am greatful to my collaborators G.Brown, J.Casalderrey-Solana, D.Teaney and I.Zahed. This work was partially supported by the US-DOE grants DE-FG02-88ER40388 and DE-FG03-97ER4014.

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