How Does the Quark-Gluon Plasma Know the Collision Energy?

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ABSTRACT

Heavy ion collisions at the LHC facility generate a Quark-Gluon Plasma (QGP) which, for central collisions, has a higher energy density and temperature than the plasma generated in central collisions at the RHIC. But sufficiently peripheral LHC collisions give rise to plasmas which have the same energy density and temperature as the “central” RHIC plasmas. One might assume that the two versions of the QGP would have very similar properties (for example, with regard to jet quenching), but recent investigations have suggested that they do not: the plasma “knows” that the overall collision energy is different in the two cases. We argue that a gauge-gravity analysis of the situation may point the way to an explanation of this important and puzzling assertion, in terms of the effects of strong magnetic fields on the plasma viscosity.
1. Peripheral LHC vs. Central RHIC

The novel form of matter produced in heavy ion collisions, the Quark-Gluon Plasma or QGP, is currently under study at two major facilities: the RHIC, which typically collides gold nuclei at maximal centre-of-mass energies up to 200 GeV per nucleon pair, and the LHC, which (in the heavy ion runs studied particularly in ALICE) typically collides lead nuclei at maximal centre-of-mass energies around 2.76 TeV per pair, latterly upgraded to 5.02 TeV [1, 2]. On this basis, it is often said that the two facilities explore two different regimes for the QGP.

There is one case, however, in which the RHIC and the LHC do probe the same regime of temperatures and energy densities. The density of a nucleus prior to a collision is not constant along a transverse direction, since the nucleus tapers away from its central axis. This simple observation has many crucial ramifications, as was pointed out by Becattini et al. [3]; inter alia it implies that the densities and temperatures arising in peripheral collisions at sufficiently high impact parameter are arbitrarily lower than those in central collisions. It follows [4–6] that sufficiently peripheral collisions studied [7, 8] at the LHC give rise to plasmas having the same energy density and temperature as plasmas produced in central collisions at the RHIC.

This opens the way to an investigation of a fundamental question: does the plasma produced in peripheral collisions “know” about global parameters like the overall impact energy, or is it sensitive only to explicitly local parameters such as the energy density?

One way to approach such questions is through studying jet quenching (see for example [9] for a review), the effect of the plasma on the propagation of highly energetic partons produced by the collision. This is conventionally described by the parameter $\hat{q}$, the mean squared transverse momentum acquired by a hard parton per unit distance travelled. In [5, 6] this is represented by a dimensionless parameter $K$, defined by

$$K \equiv \frac{\hat{q}}{2\epsilon^{3/4}},$$

where $\epsilon$ denotes the local energy density; and it is found, in the models constructed there (see also [10]), that this quantity is larger, by a factor of up to about 3, for the RHIC plasmas than for their LHC counterparts. This is true even when one compares peripheral LHC collisions with central RHIC collisions, so that, as explained above, the resulting plasmas can have the same energy density and temperature. Thus we have a strong suggestion that the local properties of the “peripheral” plasma are, in some way, (very) sensitive to a global parameter, the centre-of-mass energy of the overall collision. This strange and potentially very important development is one aspect of what has been called [11] the JET puzzle.

Work on explanations of this puzzle has begun [12, 13]. Here we wish to argue that one crucial aspect of the situation remains to be taken into account: the simple fact that central RHIC plasmas and their peripheral LHC counterparts differ in the following sense: the latter (only) are immersed in extremely powerful magnetic fields, a fact that has recently given rise to a large literature (see for example [14]). This may be relevant, because a strong magnetic field tends to suppress momentum diffusion in directions perpendicular to it [15], and possibly because it subjects the plasma to “paramagnetic squeezing” [16],

...
altering the pressure gradients; in both cases the (kinematic) viscosity of the plasma will be affected. This in turn will affect jet quenching.

That the magnetic field will modify the $K$ parameter (in the peripheral case) in some way seems clear, but it is far from clear that the modification will be significant, since the viscosity of fluids usually varies extremely slowly with pressure. Here we wish to use a simple gauge-gravity [17] model to assess whether the magnetic fields in this case, huge though they may be, can give rise to changes in $K$ of the same order of magnitude as those reported in [5, 6].

Gauge-gravity models of jet quenching are of course well known [18, 19], and such models also exist which take into account the effect of the magnetic field [20]. Here we will use a much more basic “minimal” model [21, 22] in which the bulk geometry is that of a magnetically charged dilatonic asymptotically AdS black hole. We find that this model predicts that, for plasma temperatures corresponding to central RHIC collisions and to suitably peripheral LHC collisions (with the associated magnetic field estimated as in [23]), the LHC $K$ parameter is predicted to be about 3.3 times smaller than the corresponding RHIC value; that is, it is reduced by about the same factor discussed in [5, 6].

We do not claim that this observation “explains” the puzzling sensitivity of jet quenching to the overall collision energy; holography is not so precise an instrument as that, particularly in the case of a model as simple as the one we use here; furthermore, as the authors of [5, 6] carefully point out, the magnitude of the effect itself is not firmly established. We do however wish to argue that the gauge-gravity duality suggests that a full solution of the puzzle cannot ignore the effect of the magnetic field on the plasma viscosity.

We begin by briefly outlining the “minimal” gauge-gravity model of this situation; then we state the results of our (numerical) investigation of the resulting equations.

2. A Gauge-Gravity Model of the QGP in a Strong Magnetic Field

We take the bulk geometry to be that of an asymptotically AdS dilatonic Reissner-Nordström black hole with a flat event horizon (indicated by a zero superscript), with magnetic charge parameter $P^*$ and mass parameter $M^*$: it takes the form [24]

$$g(\text{AdSdil}P^*\text{RN}^0) = -U(r)dt^2 + \frac{dr^2}{U(r)} + [f(r)]^2 \left[d\psi^2 + d\zeta^2\right],$$

where $t$ and $r$ are as usual and $\psi$ and $\zeta$ are dimensionless planar bulk coordinates corresponding to the familiar $(x, z)$ coordinates in the reaction plane at infinity. The metric coefficients are

$$U(r) = -\frac{8\pi M^*}{r} \left[1 - \frac{(1 + \alpha^2)P^*r^2}{2M^*r}\right]^{\frac{1+\alpha^2}{1+\alpha^2}} + \frac{r^2}{L^2} \left[1 - \frac{(1 + \alpha^2)P^*r^2}{2M^*r}\right]^{\frac{2+\alpha^2}{1+\alpha^2}},$$

and

$$f(r)^2 = r^2 \left(1 - \frac{(1 + \alpha^2)P^*r^2}{2M^*r}\right)^{\frac{2+\alpha^2}{1+\alpha^2}};$$

here $\alpha$ is the coupling of the dilaton $\varphi$ to the magnetic field, the corresponding term in the Lagrangian being $e^{-2\alpha\varphi}F^2$; the mass per unit event horizon area is $M^*/(\ell_P^2 f(r_h)^2)$,

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where \( r = r_h \) at the event horizon, and \( \ell_P \) is the bulk Planck length; \( P^*/(\ell_P f(r_h)^2) \) is the magnetic charge per unit event horizon area; and \( L \) is the AdS curvature scale. (This \( L \) also determines the spatial length scale at infinity, so it should reflect a typical such scale for the problem at hand.) The dilaton must be present, when magnetic fields are extremely strong (relative to the squared temperature), in order to ensure that the bulk system is mathematically consistent within string theory \([25, 21, 22]\); these considerations fix the value of \( \alpha \) when the temperature and magnetic charge of the black hole are given; otherwise the dilaton plays no direct role in the sequel.

The conformal transformation used in the gauge-gravity duality to relate the bulk and boundary geometries has the effect of imprinting the bulk magnetic field on conformal infinity: the magnetic field there is given by

\[
B_\infty = \frac{P^*}{L^3}.
\]

The equation for \( r_h \) is just, from (3),

\[
- \frac{8\pi M^*}{r_h} \left[ 1 - \frac{(1 + \alpha^2)P^{*2}}{2M^*r_h} \right]^{\frac{1}{1+\alpha^2}} + \frac{r_h^2}{L^2} \left[ 1 - \frac{(1 + \alpha^2)P^{*2}}{2M^*r_h} \right]^{\frac{2\alpha^2}{1+\alpha^2}} = 0,
\]

and the Hawking temperature of this black hole, corresponding to the temperature at infinity, is

\[
4\pi T_\infty = \frac{8\pi M^*}{r_h^2} \left( 1 - \frac{(1 + \alpha^2)P^{*2}}{2M^*r_h} \right)^{\frac{1}{1+\alpha^2}} - \frac{4\pi (1 - \alpha^2)P^{*2}}{r_h^3} \left( 1 - \frac{(1 + \alpha^2)P^{*2}}{2M^*r_h} \right)^{\frac{2\alpha^2}{1+\alpha^2}}
\]

\[
+ \frac{2r_h}{L^2} \left( 1 - \frac{(1 + \alpha^2)P^{*2}}{2M^*r_h} \right)^{\frac{2\alpha^2}{1+\alpha^2}} + \frac{\alpha^2 P^{*2}}{M^*L^2} \left( 1 - \frac{(1 + \alpha^2)P^{*2}}{2M^*r_h} \right)^{\frac{2\alpha^2}{1+\alpha^2}}.
\]

Given \( B_\infty, T_\infty \) and \( L \) (and therefore \( \alpha \)), one can use the three equations (5), (6), and (7) to solve (in physical cases) for \( r_h, P^* \), and \( M^* \), and in this way the “known” boundary parameters fix the bulk geometry. In particular, therefore, they determine (together with the bulk Planck length) the black hole entropy per unit horizon area and the black hole mass per unit horizon area. The former is \( 1/4\ell_P^2 \) for Einstein gravity, which is all we use in this minimal model; the latter is given, as above, by \( M^*/(\ell_P^2 f(r_h)^2) \). The ratio of these quantities is dual to the ratio of the entropy density of the boundary system, \( s \), to its energy density \( \epsilon \), so we have

\[
\frac{s}{\epsilon} = \frac{f(r_h)^2}{4M^*} = \frac{r_h^2 \left( 1 - \frac{(1 + \alpha^2)P^{*2}}{2M^*r_h} \right)^{\frac{2\alpha^2}{1+\alpha^2}}}{4M^*}.
\]

Given \( B_\infty, T_\infty \), and \( \epsilon \) for the boundary theory, we can use this to make a holographic prediction regarding \( s/\epsilon \).

This is relevant here because the quenching parameter \( \hat{q} \) is closely related to the entropy density. In fact \([18, 19]\), in a strictly conformal, strongly coupled boundary theory, one expects \( \hat{q} \) to scale with the square root of \( s \). (Of course, the real plasma does not correspond to a conformal theory; the errors thus inevitably introduced can however...
be estimated [18, 19], and it appears that they will not invalidate the kind of order-of-magnitude estimate we are aiming for here.)

We are now in a position to compare $K^C_{RHIC}$, the $K$-parameter for central RHIC collisions, with $K^P_{LHC}$, the value for peripheral LHC collisions resulting in plasmas at the same temperature.

3. A Holographic Computation of $K^C_{RHIC}/K^P_{LHC}(\sqrt{s_{NN}} = 2.76 \text{ TeV})$

We are interested in comparing the plasmas produced in central Au-Au RHIC collisions (at a temperature [26] of around 220 MeV after equilibration) with those produced in Pb-Pb LHC collisions (at a centre-of-mass collision energy around 2.76 TeV per pair) which are sufficiently peripheral as to give rise to the same temperature and energy density.

For the central RHIC collisions we have $B = \alpha = 0$; it is now easy to solve (6) and (7) for $r_h$ and $M^*$, and then $s^C_{RHIC}/\epsilon$ can be found from equation (8).

The energy density for central LHC collisions is estimated [27] to be around 2.3 times as large as in central RHIC collisions; using this, and following the discussion of the “thickness function” for nuclei given in [3], we find that the LHC plasmas in which we are interested arise when the impact parameter of the collision is $b \approx 12$ fm (with an assumed nuclear radius of around 7 fm). Consulting [23] one finds that this corresponds to a magnetic field of about $eB \approx 60 \, m^2/\pi$, where $m_\pi$ is the standard pion mass. These data allow us to compute $\alpha \approx 0.34$, after the manner of [21, 22]. We can now solve (5), (6), and (7) for $P^*$, $r_h$, and $M^*$, and again equation (8) yields a value for $s^P_{LHC}/\epsilon$; note that, by construction of course, $\epsilon$ is the same in both computations.

We can now proceed in a straightforward way: we have, in an obvious notation,

$$\frac{K^C_{RHIC}}{K^P_{LHC}} = \frac{\hat{q}^C_{RHIC}/2\epsilon^{3/4}}{\hat{q}^P_{LHC}/2\epsilon^{3/4}} = \frac{\hat{q}^C_{RHIC}}{\hat{q}^P_{LHC}} = \sqrt{\frac{s^C_{RHIC}}{s^P_{LHC}}} = \sqrt{\frac{s^C_{RHIC}/\epsilon}{s^P_{LHC}/\epsilon}},$$

and this last quantity is something we can, as explained, compute holographically. The final result is as follows:

$$\frac{K^C_{RHIC}}{K^P_{LHC}(\sqrt{s_{NN}} = 2.76 \text{ TeV})} \approx 3.34.$$  

This is indeed of the same order of magnitude as the result reported in [5,6].

We conclude that the holographic approach indicates that the local plasma in the peripheral case “knows” about the global impact energy through the magnetic field. On a qualitative level, this is not entirely unexpected. To understand why, note that the relevant property of the plasma here is its “momentum diffusivity” or kinematic viscosity $\nu$, defined as the ratio of the dynamic viscosity $\eta$ to the energy density. Then we have

$$\nu = \frac{\eta}{\epsilon} = \frac{\eta}{s} \times \frac{s}{\epsilon} = \frac{1}{4\pi} \times \frac{s}{\epsilon},$$

where we have used the well-known KSS relation [28] in the last step. From this we see that the kinematic viscosity is, according to holography, affected by the magnetic field in just the same way as $s/\epsilon$: that is, these extremely intense fields reduce the momentum diffusion to a significant extent. This can be expected to influence the jet quenching parameter, due to the effects mentioned earlier.
4. A Holographic Prediction for $K_{RHIC}^C/K_{LHC}^P(\sqrt{s_{NN}} = 5.02 \, TeV)$

The calculations in [5,6] are primarily concerned with LHC collisions at 2.76 TeV per pair. The holographic technique is easily extended to the recent runs [29] at 5.02 TeV per pair, as follows.

The energy density for central collisions at very high collision energies is discussed, for example, in [30]. We will assume that a typical energy density for central Pb-Pb LHC collisions at 5.02 TeV per pair is around 3 times the maximal RHIC density. Using this in the same manner as above, we find that the impact parameter of the relevant collisions (resulting in a plasma with the same temperature, about 220 MeV, as plasmas produced in central RHIC collisions) is slightly larger than before, $b \approx 13$ fm. Again consulting [23] (noting in particular that the magnetic field at a given location increases roughly linearly with the impact energy), one finds that now $\epsilon B \approx 120 \, m_e^2/\pi$. We compute $\alpha \approx 0.36$ with these data. Solving (5), (6), and (7) in this case and using (8), we compute a value for $s_{LHC}^P/\epsilon$ which is still smaller than the value for 2.76 TeV collisions, and the final result is as follows:

$$\frac{K_{RHIC}^C}{K_{LHC}^P(\sqrt{s_{NN}} = 5.02 \, TeV)} \approx 4.15;$$

that is, holography predicts a roughly 25% increase in this quantity as we move from 2.76 TeV to 5.02 TeV collisions. The corresponding estimate in [5] is for an increase of around 15%; again we remind the reader that the gauge-gravity technique is not capable of precision such as to render this difference significant. One should note also that, at such high values of the impact parameter, various other factors we have neglected (such as the fact that the nuclei are not spherical) will come into play. It will be interesting to study this further as the data are analysed more completely.

5. Is the Effect Independent of Centrality?

In [3,6] it is claimed that the value of the $K$-parameter in collisions at a given impact energy is independent of the centrality of the collision, at least for transverse momenta that are not very small (note in this connection that holography may not be reliable [21] for LHC collisions at small impact parameters, that is, for ultra-high densities). This presents a challenge for the “magnetic” theory, because the magnetic fields are much smaller at low values of the impact parameter than at high values.

We have computed (in the same manner as above) the ratio of the holographically predicted value of $K$ for the LHC plasmas produced in peripheral collisions at 2.76 TeV per pair, as a function of the impact parameter $b$ (denoted $K_{LHC}^P(b)$), to its value for central LHC collisions (denoted $K_{LHC}^C$), also at 2.76 TeV. The results are shown in Figure 1, with $b$ in units of fm. We see that, beyond about $b = 3$ fm, the graph is indeed remarkably flat. It appears that the magnetic field has an almost immediate effect in reducing $K_{LHC}^P(b)/K_{LHC}^C$, even at relatively small values of the field; the subsequent large increases in the field have almost no further effect. (In fact there appears to be a broad local minimum around $b \approx 9$ fm.) Thus we can claim that the simple gauge-gravity model of the magnetic effect is also in accord with this finding of [3,6].
6. Conclusion: The Magnetic Field is the Key

The puzzling claims \[5, 6\] that jet quenching can detect the difference between the QGP produced in central RHIC collisions and the plasma produced in peripheral LHC collisions, even when the local temperature and energy densities are the same, and that this effect is independent of the centrality, demand an explanation. We have argued, using a very simple gauge-gravity model, that the extremely intense magnetic fields arising in one case (the LHC plasmas), and not the other, may be the key to understanding these claims. The model suggests \[a\] that the magnetic field has an unexpectedly strong effect on jet quenching, possibly by reducing the momentum diffusivity of the plasma, and that \[b\] this effect arises immediately at relatively low fields and changes little with further increases associated with larger impact parameters, as seen in Figure 1.

The conclusion is simply that further investigations of these intriguing observations should focus on the effect of the magnetic fields arising in peripheral heavy-ion collisions. In particular it would be useful to have a better understanding of the effect of strong fields on QGP viscosity, focusing perhaps on paramagnetic squeezing \[10\], and using more sophisticated gauge-gravity models than the one employed here.

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