Deconfinement, Monopoles and New Phenomena in Heavy Ion Collisions

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We discuss various manifestations of the “magnetic scenario” for the quark-gluon plasma viewed as a mixture of two plasmas, of electrically (quark and gluons) as well as magnetically charged quasiparticles. Near the deconfinement phase transition, \( T \approx T_c \) very small density of free quarks should lead to negligible screening of electric field while magnetic screening remains strong. The consequence of this should be existence of a “corona” of the QGP, in a way similar to that of the Sun, in which electric fields influence propagation of perturbations and even form metastable flux tubes. The natural tool for its description is (dual) magnetohydrodynamics: among observable consequences is splitting of sound into two modes, with larger and smaller velocity. The latter can be zero, hinting for formation of pressure-stabilized flux tubes. Remarkably, recent experimental discoveries at RHIC show effects similar to expected for “corona structures”. In dihadron correlation function with large-\( p_t \) trigger there are a “cone” and a “hard ridge”, while the so called “soft ridge” is a similar structure seen without hard trigger. They seem to be remnants of flux tubes, which – contrary to naive expectations – seem to break less often in near-\( T_c \) matter than do confining strings in vacuum.

Keywords: confinement, monopoles, quark-gluon plasma, flux tubes

1. Introduction

Let me start my written version of the talk in the same way as the oral one: with comments on my interaction with Misha Shifman. It is more than 30 years since he starts visiting Novosibirsk, working with Arkady on QCD sum rules, while I occasionally managed to visit ITEP. 1970’s were formative years for all of us, and I learned a lot from those discussions. While we shared many interests over these years – to QCD correlators, heavy quark hadrons, instantons, now monopoles and confinement – we have so different styles that we never actually worked on exactly the same thing.
That produced some distance, a different perspective and often mutual criticism. And yet, it only strengthened our friendship: we are in agreement on so many other things in life. So, it is a delight to see Misha at another conference, of his beloved “Continuous Advances” series in particular, full of life and as usual having a joke or a couple of good (if sceptical) questions to any talk.

This talk fits into this long-time pattern perfectly. While in the years 2000-2005 I was preoccupied by RHIC data and then various strongly coupled plasmas, AdS/CFT and other issues not close to Misha, my talk at Continuous Advances 06 included a part about “magnetic scenario” for temperatures near $T_c$, on a plasma side. Presently both of us study different aspects of the “magnetic side” of QCD, trying to understand confinement mechanism: while Misha looks at models with enough supersymmetry to have rigid theoretical control, I looks at RHIC and lattice data plus some generic theoretical arguments to get the picture. Recent results and experimental discoveries look quite exciting: I hope that makes a good addition to this proceeding, a kind of our collective birthday present.

2. Dual electric/magnetic plasmas and transport

Let me thus start by reminding what I was speaking about at CA06. Among very few things we can agree on with Misha is a general “dual superconductor” view of confinement by ’t Hooft and Mandelsham. If so, the confinement transition must be Bose-Einstein condensation (BEC) of certain magnetically charged objects. Furthermore, right above $T_c$ there should be sufficient density of them in a “normal” (uncondensed) state.

Dirac’s quantization condition demands the product of electric and magnetic couplings to be an integer: thus those two couplings must run in the opposite directions (as it happens e.g. in Seiberg-Witten $\mathcal{N}=2$ theory). So Liao and myself proposed a view of CD plasma as of dual mixture, with two Coulomb-interacting subsystems, both at strong coupling whose magnitude is running with $T$ in the opposite directions. (Chernodub and Zakharov have independently proposed a similar picture, with large role of monopoles in the near-$T_c$ region.) But how to see if this is true?

On the lattice one can identify monopoles by “the end of the Dirac string” in a certain gauge: so one can look at monopole-antimonopole correlations in space. An example for two temperatures is shown in Fig.1

\[\text{Let me note that thus what is called a monopole actually include dyons, and that small size of the core and coarse lattices does not allowed for detail knowledge of its structure.}\]
(left). Note characteristic peaks which are found in strongly correlated liquids. Note also that the higher $T$ has stronger correlation: this implies that magnetic Coulomb coupling does grow with $T$, contrary to asymptotic freedom for electric coupling. The right figure is our molecular dynamics (MD) simulations for plasmas containing monopoles: note that the red one (monopole-antimonopole) has the same shape and magnitude as the lattice data on the left. So, a picture two dual plasmas, electric and magnetic ones, does work.

What is the contribution of the monopoles to transport? Before we used classical molecular dynamics while this time I report results by Ratti and myself. We solved a beautiful quantum-mechanical problem of gluon-monopole scattering. Asymptotically, at $T \to \infty$, this process is subleading by one power of $\log(T)$. Yet at $T$ few times $T_c$, when the monopole density is still much smaller than that of quarks and gluons, we found that gluon-monopole scattering actually dominates the transport. This happens because of characteristic large backward scattering, absent in charge-charge case. Scattering rate and $(\eta/s)$ are shown in Fig. 2(left), together with the ones obtained for the usual perturbative gluon-gluon scattering. There is a qualitative agreement between these results and the experimental value for $\eta/s$ observed at RHIC (the box on the left), as well as with the famous AdS/CFT result $\eta/s = 1/4\pi$. LHC will create (in PbPb collisions) QGP with $T \sim 4T_c$: we thus have our prediction of how good a liquid QGP will
be there.

3. Universal Feynman’s criterion for Bose-Einstein condensation and monopoles at deconfinement

If one wanders whether confinement is indeed a BEC of monopoles, looking at their behavior on the lattice close to $T_c$ is a good idea. But before doing it, one should know what to look at. In a paper with Marco Cristoforetti we had revived an old idea by Feynman, in which he introduced the notion of the universal critical value for the per-particle action in large supercurrent clusters. Note that in Matsubara formalism bosons should make periodic paths: but it may be not for each particle separately but for say k particles, called k-cluster. Feynman argued that the sum over cluster size k should diverge at the critical point, when action suppressing “jumps” from one particle location to the next is balanced by multiplicity of clusters. Feynman included kinetic energy only and thus approximated it by the straight line, so his expression for the amplitude is

$$y_F = \exp \left[ - \frac{m^* T d^2}{2 \hbar^2} \right]$$  \hspace{1cm} (1)

Marco and I decided to work it out more, and include also potential energy. We had built the “moving string” model for the supercurrent, a line of particles in a crystal (a reasonably good model for liquid He) and calculated the action of all the string moving as a whole using realistic atomic potentials. Semiclassical calculation (another use of periodic instantons or calorons) or numerical path integration have given close actions. In
Fig. 2(right) we plot the ratio of the calculated critical temperature $T_c/T_0$, where $T_0$ is Einstein’s critical temperature for free Bose gas. We found that for strongly interacting liquid He the critical action is $S_c = 1.655$, which is exactly the same value as one would get for noninteracting gas. So, Feynman was right about universality of his criterion: he just have not done a complete calculation to verify it for liquid He!

Now, with confidence about universal BEC criterion, we applied it for QCD monopoles. Since we know approximately their density, we can set upper limits on the monopole mass and magnetic Coulomb coupling at $T_c$. It turns out they must be as light as about 200 MeV or so.

Does it all agree with the lattice behavior of monopoles? As we were told by D’Elia, clusters with $k=1,2,3$ indeed have equal probability at $T_c$. The monopole mass is not yet accurately measured, but it indeed strongly decreases toward $T_c$ about as low as predicted. In summary: monopoles do behave similarly to He atoms at their respective BEC transition temperatures.

4. QGP corona: theory

The word “corona” used here comes from the physics of the Sun. Let me briefly remind that it was started by Galileo Galilei, who in 1612 spent some time observing the motion of the black spots on the Sun and correctly concluded from motion of the spots that they must resign on a surface of a rotating sphere: he thus argued the spots were not shadows of some planets passing in front, as it was thought of before. In due time relation between the spots and solar magnetism was understood: modern telescopes allows one to see the fine structure of solar spots, resolving individual magnetic flux tubes. Better understanding of solar magnetism came with the advance of plasma physics in 1940’s and development of MHD, which explained both the influence of diffuse magnetic field on plasma and formation and mechanical stability of the flux tubes.

Early stages of heavy ion collisions are believed to be described by the so called “glasma”, a set of random color fields created by color charges of partons of the two colliding nuclei at the moment of the collision. However as two discs with charges move away from each other, those classical field are getting smaller and (in a still poorly understood process) rather quickly create the quark-gluon plasma, in which the occupation numbers are thermal, $O(1)$, with the classical fields disappearing.

Transition from field to plasma has been subject of multiple works, too many to be mentioned. One notable idea is existence of instabilities such
Fig. 3. (left) A snapshot of unscreened electric (dual-magnetic) field in the M (near-$T_c$) region of the fireball. Upper and lower ones correspond to full RHIC energy and the reduced energy (analogous to SPS). (right) A sketch of the transverse plane of the colliding system: the “spots” of extra density (a) are shown as black disks, to be moved by collective radial flow (arrows). Naïve sound expansion (b) would produce large-size and small amplitude wave: yet the correct solution includes also brighter secondary wave (c) of smaller radius.

as the so called Weibel instability, producing filamentation of the longitudinal flow and transverse magnetic field $B_\perp$. Asakawa et al.\textsuperscript{12} argued that such chaotic magnetic field would remain in plasma, leading to “abnormal viscosity”. Unfortunately those ideas, derived perturbatively and basically taken from electromagnetic plasma contexts, cannot possibly be valid for QGP at RHIC.

The crucial difference between the QED and QCD plasmas lies in the existence of magnetically charged quasiparticles – monopoles and dyons – leading to nonzero magnetic screening mass $M_M \sim g^2 T$ first suggested by Polyakov long ago\textsuperscript{13} and by now well confirmed by subsequent lattice studies. Thus hot QGP, unlike the electromagnetic plasmas, screen both the electric and the magnetic fields at some microscopic scales, although at a bit different ones. As at RHIC the central part of the produced fireball reaches relatively high temperature $T \sim 2 T_c$, we expect both $E, B$ fields to be effectively screened there, see the central cylindrical part of Fig.3(upper
left) marked QGP. But in the near-$T_c$ region there is magnetic plasma in
the outer cylindrical part of Fig.3 marked M (mixed or magnetic). It is very
important to emphasize that although this region on the phase diagram
is represented by a very narrow strip $|T - T_c| \ll T_c$, it corresponds to
more than order of variation of the energy or entropy density, and the
corresponding space-time volume in the expansion of the fireball is by no
means small. Another snapshot of the geometry of the M region is shown
in Fig.3(lower left): it corresponds to lower collision energy for which the
M-phase resign inside the fireball: obviously this case is very different, it is
planned to be investigated in a specialized RHIC run in 2011.

Multiple lattice studies, e.g. ref. have shown that at $T < 1.4 T_c$ the rela-
tion between electric and magnetic screening masses get inverted, namely
$M_M > M_E$. As $T \to T_c$ the electric screening mass strongly decreases,
partly because of heavy quark and gluon quasiparticles and partly because
of their suppression by small Polyakov loop expectation value $< L >$. So,
near $T_c$ $M_M(T \to T_c) \approx 3 T_c, M_E(T \to T_c) \approx 0$. The most important
consequence is that in the M-phase of the collision there is dense mag-
netic plasma, in which magnetic field is well screened while electric one
remains unscreened for a significant time. This leads to the central idea:
that QGP should have a “dual corona” in which electric (rather than mag-
netic) fields coexist with the plasma, affecting both the overall expansion
and the propagation of perturbations. We will suggest to use “dual magne-
thydrodynamics” (DMHD) for the description of diffuse electric fields in
the M-phase, in particular study their effect on the velocity of propagation
of small perturbations. We will further argue below that like solar corona,
that of QGP should have metastable flux tubes, although microscopically
thin ones which cannot be directly described by DMHD approximation.

Magnetohydrodynamics is a well known part of plasma physics, devel-
opled by Alfven, Fermi, Chandrasekhar and many others. It is an approxi-
mation which keeps only one (magnetic) field nonzero and half of Maxwell
eqns, while the other (electric) field is assumed to be totally screened. Ideal
MHD approximation is the limit of infinite conductivity of plasma $\sigma \to \infty$,
similar to zero viscosity approximation for ideal hydrodynamics. In MHD
the coupling between the field and and matter is obtained by inclusion of
the (magnetic) field contribution into the stress tensor of the medium.

Due to space limitations I will not put a complete set of resulting
equations here, jumping directly to the main consequences. In a
“(dual)magnetized” plasma small amplitude perturbations are split into
two modes, known as Alfven waves, propagating with two different speeds

\[ u_\pm^2 = \left( \frac{1}{2} \right) \left[ \frac{E_2^2}{\epsilon} + c_s^2 \right] \pm \left( \frac{1}{2} \right) \left[ \frac{E_2^2}{\epsilon} + c_s^2 \right]^2 - \frac{4E_2^2 \cos^2 \theta c_s^2}{\epsilon} \right]^{1/2} \]  

where \( E, c_s, \epsilon \) are the electric field, the speed of sound in “unmagnetized” medium without field and plasma energy density, \( \theta \) is the angle between the field strength and the direction of the wave propagation. Note a case in which the wave goes transverse to the field (\( \cos \theta = 0 \)) in which the lower mode has zero speed. In the case of jet quenching – when the jet direction is more or less up to experimentalist to pick – general shape of these waves can be complicated. However when the jet (or original charge fluctuation) propagates longitudinally, in the same direction as the field, the problem is axially symmetric and results in general in two cones. The angles of their propagation can be obtained from the previous expression, in which the l.h.s. is substituted by Mach relation \( u \rightarrow \cos(\theta)v \), \( v \) is the velocity of the jet, and solve it for the \( \cos \theta \) (see the paper for more details on solutions).

If instead the electric field to be constant in space there is a a spot, localized in transverse plane, inevitable there is nonzero \( \partial E_z / \partial r \), which is a part of \( \text{curl}(\vec{E}) \) and by the (dual) Maxwell equation it is proportional to (dual, or magnetic) current \( \tilde{j}_0 \). This tells us that a flux tube solution must have a “coil” with a current running around and cancelling the field outside the spot. Ideal DMHD has simple axially symmetric solution – the macroscopic flux tube, with the field pressure balancing that of the plasma. We dont discuss them in detail here because we think the flux tubes appeared in QGP corona are actually microscopically thin, and so should be considered in a way worked out by Liao and myself previously.\(^7,10\) Let me emphasize it again: we don’t speak hear about well known flux tubes at zero or low \( T \) but about flux tubes which exist at \( T > T_c \), in the deconfined phase with magnetic constituents. As the current is not a supercurrent in this case, they are not stable forever but are metastable. We worked out quantum mechanics of the monopole-flux tube scattering and conditions for its mechanical stability. Perhaps surprising to many, those flux tubes are tighter (have larger tension) than the vacuum ones: this comes because uncondenced monopoles have thermal momenta and their scattering creates larger pressure on the tube. As we will discuss at the end of the talk, those tubes in the M-phase seem to have smaller breaking rate as well: so they are tighter in any sense.

The original hints for their existence came from lattice data on finite-\( T \)
potentials between static quark and antiquark, studied on the lattice. In brief, the central observation is large difference found between the free energy $F(T, r)$ and the potential energy $V(T, r) = F(T, r) + TS(T, r)$ associated with quark pair at distance $r$. While $F$ has no linear part (in $r$) above $T_c$, both $V$ and $S$ have it. The physical difference between the two is that $F$ corresponds to adiabatically slow motion of the quarks, slow enough to produce maximal entropy possible and reach thermal equilibrium at any $r$. However when quarks are moving with certain time scale, only a fraction $x$ of the maximal entropy can be produced (because of Landau-Zener argument on level crossing): thus the effective potential is $V_{\text{eff}}(T, r) = F(T, r) + (1-x)TS(T, r)$. For relatively rapid motion in which no entropy is produced, $x = 0$, and one returns to $V(T, r)$. The difference between the two potentials can be also explained in other way: $F(T, r)$ is related to a stable flux tube with a “supercurrent coil” while $V(T, r)$ has just “normal” coil and thus is metastable.

5. QGP corona: experiment

Three different correlation phenomena have been discovered in heavy ion collisions at RHIC:

(i) The “cone” has been discovered in the 2-particle azimuthal correlations. The “away-side” peak at $\Delta\phi = \pi$ due to a partner jet, balancing hard trigger particle, disappears, substituted by new peaks at completely different angle. After discovery of those effects there was extensive studies of the 3-particle correlations, which confirmed that the observed structure is indeed cone-like, and not e.g. a reflected jet.

(ii) the hard ridge is also seen in 2-particle correlators, but plotted on the two-dimensional $\Delta\phi - \Delta\eta$ plane, the differences between the azimuthal angles and pseudorapidities of the two particles. The jet remnants make a peak near $\Delta\phi = 0, \Delta\eta = 0$, which was found to sit on top of the “ridge”, with comparable width in $\Delta\phi$ but stretched in rapidity to nearly all 10 units (between beam rapidities at RHIC). For plots and various features one can consult the original talk by Putschke.

(iii) the “soft ridge” is found by STAR collaboration without a trigger, in the 2-particle correlations. For experimental details and phenomenological considerations see talks at recent specialized workshop.

We will return to these observations below, turning now to their suggested explanations:

(i) Stoecker et al, as well as Casalderrey, Teaney and myself have proposed
that the energy deposited by a quenched jet goes into two hydrodynamical excitation modes, the sound and the so called diffusion or wake modes. The sound from the propagating jet should thus create the famous Mach cone, in qualitative agreement with the conical structure observed. However surprising experimental fact is quite large amplitude of this signal, as well as large value of the cone angle which has not been reproduced by hydrodynamics (or AdS/CFT). The cone angle is in the range $\theta_{\text{cone}} = 1.2 \sim 1.4$ radians (not too far from $\pi/2 = 90^\circ$ or cylindrical waves!) which by Mach formula gives the speed of pertinent perturbation to be about

$$< v_{\text{wave}} > = \cos(\theta_M) \approx 0.2 \quad (3)$$

while the expected speed of sound is .3 at its lowest point near $T_c$ and about .5 in QGP.

(ii) One early model for “hard ridge” has been introduced in my paper.\textsuperscript{20} It relates it with the forward-backward jets accompanying any hard scattering, providing extra particles (“hot spot”) widely distributed in rapidity. This idea is then combined with the one suggested previously by Voloshin,\textsuperscript{21} namely that extra particles deposited in the fireball would be moved transversely by the radial hydrodynamical flow, should produce a peak at certain azimuthal angle corresponding to the position of the hot spot, see Fig.3(upper right). While particles of the ridge are separated by large rapidity gaps and cannot communicate during the expansion process, their azimuthal emission angles remain correlated with each other because they originate from the same “hot spot” in the transverse plane.

(iii) Similarly, transverse hydro boost of “hot spots” was used for the explanation of the “soft ridge” by McLerran and collaborators.\textsuperscript{22,23} They have pointed out that the initial state color fluctuations in the colliding nuclei would create “spots” at some positions in the transverse plane. The idea is shown in Fig.3(upper right): the spots can be carried by hydro expansion and get visible at certain azimuthal angles. Further confirmation of hydro origin of ridges comes from the centrality dependence of the angular width of the ridge: the peak in azimuth sharpens for more central collisions.

So, at a very qualitative level the origin of all three phenomena seem to be explained: yet at more qualitative level a lot of puzzles appear. As an example, consider the simplest of them, the “soft ridge”. As discussed in,\textsuperscript{22,23} the initial stage (proper time $\tau \sim 1/Q_s \sim 0.2 \, fm/c$ where $Q_s \sim 1 \, GeV$ is the so called saturation scale at RHIC) can be discussed using classical Yang-Mills equations: thus color fluctuations naturally appear. However, the observed pions come from final freezeout time, separated from the ini-
tial “glasma” era by much longer time $\tau \sim 10 \, \text{fm}$. This is certainly so, as the explanation heavily relies on radial hydro velocity and thus it has to wait till the hydro velocity is being created. As we will argue below, there are many reasons why one might have expected nearly complete disappearance of this signal during this time.

Common to all three cases is deposition of some additional energy (or entropy), on top of the “ambient matter”. The number of correlated particles in all of them constitute a small ($\sim 10^{-3}$) fraction of the total multiplicity: thus they can only be seen in a high-statistics correlation analysis.

Similar to circles from a stone thrown into a pond, initial perturbation should become some expanding waves, and hydrodynamics predicts that basically nothing is left at the original location at later times. Even ignoring any dissipation and using ideal hydrodynamics one finds the radius of those waves at final freezeout to be given by the “sound horizon”

$$ R_h = \int_0^{\tau_f} d\tau c_s(\tau) $$

So, by the freezeout proper time $\tau_f \sim 10^{-15} \, \text{fm/c}$, this distance is large, $\sim 6 \, \text{fm}$ or so, see Fig.3(right middle). The amplitude of the sound wave is decreasing accordingly, and the width of $\phi$ distribution grows. Simple calculation show that in such case the width of the peak would be larger than observed, and the amplitude smaller. This is a common puzzle for all three phenomena mentioned: hydro fails to explain why all of them remain observable, as if nothing happened to them during rather long time of the hydro process. The second set of waves provided by DMHD, with smaller velocity of propagation, or better yet stabilized flux tubes is thus suggested as an explanation, see Fig.3(bottom right).

Let me end this brief wandering into experiment-generated puzzles by mentioning recent RHIC data which provided direct experimental indications for enhanced stability of the flux tube in matter relative to pp. Those are from PHOBOS collaboration (now with a detector no longer physically existing) which has large rapidity coverage of their silicon detector. They have studied clustering of secondaries in rapidity and found that the number of charged particles in a cluster observed in AuAu collisions is about twice that seen in correlation studies of the pp collisions, reaching the size of about 10 particles. The rapidity width of the cluster is also larger than for those in pp: they are not consistent with isotropically decaying resonances but apparently elongated longitudinally. We basically see fragments of the flux tubes, in pp and AuAu collisions: and apparently the tubes decay less frequently in the latter case, into larger pieces. I take it as the most direct
indication to existence of the “magnetic plasma” from the data.

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