Magnetic component of gluon plasma and its viscosity

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We discuss the role of the magnetic degrees of freedom of the gluon plasma in its viscosity. The main assumption is that motions of the magnetic component and of the rest of the plasma can be considered as independent. The magnetic component in the deconfined phase is described by a three-dimensional (Euclidean) field theory. The parameters of the theory can be estimated phenomenologically, from the lattice data. It is not ruled out that the magnetic component is superfluid.

I. INTRODUCTION

The interpretation of heavy ion experiments at RHIC and first-principle lattice simulations suggest that the quark-gluon plasma has quite unusual properties [1, 2]. Contrary to general expectations, at temperatures just above the critical temperature, $T_c$, the plasma has properties rather of an ideal fluid than of a weakly interacting gas of quarks and gluons [3].

The unexpectedly low viscosity is in contrast with thermodynamical properties of the plasma which do not betray much unexpected. Indeed, at high enough temperatures the difference between the observed pressure (energy) density and its perturbative value can be fitted by a $g^6(T)$ contribution,

$$p(T)_{\text{full}} - p(T)_{\text{pert}} \approx \text{const} \cdot g^6(T),$$

where $g^2(T)$ is the running constant. The fit (1) is expected on theoretical grounds since in the order $g^6(T)$ one runs into infrared divergencies which can in fact be treated only non-perturbatively. Thus the constant in the r.h.s. of Eq. (1) is not calculable analytically at this time.

Let us consider an idealized picture and assume for the moment that the deviations of the thermodynamical properties from (weak-coupling) perturbative values are small while the viscosity is much lower than its perturbative value. This would suggest that we deal with two independent fluid components with viscosities $\eta_1, \eta_2$ then we have additivity of fluidity, i.e., inverse viscosity $\eta$:

$$\frac{1}{\eta_{\text{tot}}} = \frac{(\text{phase space})_1}{\eta_1} + \frac{(\text{phase space})_2}{\eta_2}.$$ (2)

If, say, $\eta_2 \approx 0$ the total value $\eta_{\text{tot}}$ can still be small even if the corresponding phase space factor $(\text{phase space})_2$ is small.

In this note we will explore the possibility that the magnetic component of the Yang-Mills plasma provides us with an independent motion in the sense of viscosity, see Eq. (2). Examples of “independent motions” in condensed-matter systems are well known. In the case of ordinary superconductivity, the contribution of Cooper pairs is independent of electrons in the normal state. Closer to our problem, superfluid and ordinary (or dissipative) components of liquids at low temperatures can be treated as independent [11].

In a crude approximation, one can understand magnetic component the 3d field theory which determines the r.h.s. of Eq. (1) at high temperatures. Indeed it is known since long [12] that at high temperatures it is the 3d field theory corresponding to the zero Matsubara frequency $\omega_M = 0$ which is to be treated non-perturbatively. In this limit, the temperature dependencies of all the non-perturbative observables can be reconstructed from their dimensions. For example, the string tension of the spatial Wilson line is to be proportional to $g_3^2 = g^2(T)T$, (4)

$$\sigma_3 \sim g_3^2,$$

where

$$g_3^2 = g^2(T)T.$$ (3)

is the dimensionally reduced gauge coupling, $g^2(T)$ is the running coupling of the 4d Yang-Mills theory. Here $g^2(T)$ is assumed to be small enough to serve as a small expansion parameter. Numerically, the scaling laws like those in set in at temperatures not too much higher than the critical temperature of the deconfining phase transition $T_c$ although $g^2(T)$ does not seem yet to be small ($T_c/\Lambda_{\text{QCD}} \approx 1.2$).

More precisely, the magnetic component is defined in terms of the magnetic monopoles and center vortices identified on the lattice. These degrees of freedom are commonly believed to be responsible for confinement at low temperatures, for a review see, e.g., [13]. As is argued in Refs. [5, 6, 7, 8, 9] at $T > T_c$ the magnetic degrees of freedom become a part of the Yang-Mills plasma. A subtle point is that magnetic degrees of freedom are studied on the lattices, or in Euclidean space while viscosity is defined most straightforwardly in the Minkowski space. In particular, we are going to treat the magnetic component as an “independent motion” in the Minkowski space.
To substantiate (or reject) this hypothesis one needs a continuum-theory interpretation of the lattice defects. Dual models of Yang-Mills theories, see in particular [8] and references therein, seem to provide such an interpretation. Namely, the dual models are formulated in terms of strings living in extra dimensions, for a review see, e.g., [12]. Then there exist various topologically stable solutions in the dual formulation of the Yang-Mills theories. In particular, the observed properties of the lattice vortices and monopoles fit remarkably well the pattern expected within the dual models for the magnetic strings [8, 10]. Moreover, the monopole and vortex pictures get unified since monopoles are to be thought about as 1d defects (trajectories) living on the 2d defects (vortices, or strings) [16, 17]. The monopoles and vortices constitute the magnetic component of the gluon plasma.

What is most relevant to our purposes, it is expected theoretically that the magnetic strings become time oriented at \( T > T_c \) since only time oriented magnetic strings are (nearly) tensionless at \( T > T_c \). Then the magnetic strings reduce to their projections to a time-slice since the time dependence is trivial. Thus, the solutions of the full 4d theory are mapped into 3d solutions. Consider now the 3d medium of these topological excitations. Since we deal with solutions of the full theory we do not need to consider further, for example, their interaction with gluons. The properties of the 3d medium depend, however, on the interaction of the topological excitations between themselves which is not taken into the account yet. Similarly, in the theory of superconductivity one starts first with an (approximate) solution for the Fermi-liquid at \( T=0 \). The Cooper pairs emerge after accounting for (relatively weak) interactions near the edge of the Fermi-sphere.

At present, there are no means to clarify interaction among the solutions theoretically and we will rely on the lattice phenomenology at this point. The lattice data are in the Euclidean space, however. In the static approximation for the magnetic defects the continuation from the Euclidean to Minkowski space is trivial and this is the approximation we will use. In the static approximation, the measurements reduce to the measurements on the ground state of the 3d system (which is the same in the Euclidean and Minkowski spaces). There is a spectrum of excitations which determine, in particular the time development of the system. The spectrum is obtained by quantization on the background of the classical solutions. If there is time dependence, the continuation from the Euclidean to Minkowski space is highly non-trivial and very difficult in reality. However, understanding the ground state alone allows one to decide, for example, whether we deal with a superfluid. This is our strategy here.

**II. MAGNETIC COMPONENT OF THE PLASMA**

At high temperatures and for static quantities, all the non-perturbative physics is expected to be described in terms of a three-dimensional theory [12]:

\[
L = \frac{1}{g_3^2} \left( \frac{1}{2} \text{Tr} F_{ij}^2 + |D_i \Pi^a|^2 + V(\Pi^2) \right),
\]

(5)

where \( \Pi^a \) is a scalar color field (\( \omega_M = 0 \) component of the potential \( A_0^a \)). As is mentioned above, the dimensional reduction implies simple scaling laws for various quantities. In particular, if one defines magnetic monopoles within the 3d Yang-Mills theory (which is a part of [4]) then the monopole density should scale as \( g_3^6 \) in terms of the dimensionally reduced coupling \( 4 \). And, indeed, in the 3d Yang-Mills theory [13]:

\[
\rho_{\lambda, \text{mon}} \approx 10^{-7} g_3^6.
\]

(6)

According to [4] the density is proportional to \( T^3 \) as would be also the case for massless particles at temperature \( T \). However, the density [48] is not given by the Boltzmann distribution in terms of the original temperature. The temperature dependence arises because of the rescaling the fields of the original 4d theory. This trivial observation becomes crucial later.

As is mentioned in the Introduction, within the dual model of Yang-Mills theory there exists an absolutely different mechanism of reducing the non-perturbative physics from four to three dimensions [19]. It is related to the dynamics of strings (which are the basic objects of the dual formulation, or Yang-Mills theories in the infrared). To put it simply, instead of 2d magnetic surfaces or strings percolating in 4d at \( T = 0 \) one has at \( T \geq T_c \) particles percolating in 3d. Such a percolation can be adequately described by 3d field theories. This conclusion, as is argued in [9], arises within various approaches, such as models of gauge/string dualities [8], effective field theories, see in particular [20], or approaches based on the lattice data as referred to in [9]. For the sake of our presentation we will reiterate the main points in the language of the lattice defects, i.e. 2d surfaces (strings) or 1d trajectories (monopoles) mentioned above.

It is useful to start from the \( T = 0 \) theory of confinement. Confinement is commonly believed to be due to condensation of the magnetic degrees of freedom. Usually one understands by the magnetic degrees of freedom either Abelian monopoles or \( Z_2 \) vortices, for a review see, e.g., [14]. In fact both projections are manifestations of one and the same non-Abelian object. Phenomenologically, the monopole trajectories cover densely the vortices (2d surfaces) or, vice versa, the vortices can be defined as minimal-area surfaces spanned on the monopole trajectories, for a review see [15]. Both the vortices and monopoles percolate through the vacuum state, i.e. form infinite clusters of the 2d surfaces or 1d trajectories. Important
properties of these clusters is that their total length, respectively area, scale in physical units:

\[ L_{\text{mon}}^{\text{tot}} \approx \text{const} \cdot \Lambda_{QCD}^3 V^{(4)}_{\text{tot}}, \]
\[ A_{\text{vort}}^{\text{tot}} \approx \text{const} \cdot \Lambda_{QCD}^2 V^{(4)}_{\text{tot}}, \]

where \( V^{(4)}_{\text{tot}} \) is the total 4d volume of the lattice. The detailed picture for confinement does depend on whether one uses monopoles or vortices. The existence of two alternative languages, as we will see, is important within the context of this note.

What happens at \( T \geq T_C \) is that the 4d percolation of the defects is becoming a 3d percolation. In more detail and in the monopole language \([8,9]\) the 4d percolating cluster disappears. Which corresponds to destroying the condensate of the magnetically charged field by temperature. Instead there appear monopole trajectories which are wrapped around the periodic time direction. One can argue that the wrapped trajectories in the Euclidean space correspond to real particles in the Minkowski space \([7]\). The 3d density of the wrapped trajectories scales indeed in physical units \([7]\):

\[ \rho_{\text{wrt}} \approx T^3 f(T, \Lambda_{QCD}), \]

where the function \( f(T, \Lambda_{QCD}) \) is slowly varying at high temperatures. It is crucial however that this function does not depend on the lattice spacing, as it should be for physical objects.

Phenomenologically the geometrical picture simplifies actually further. First, already at \( T \) close (and larger than) \( T_c \) the wrapping number is equal to \( n_{\text{wrt}} = 1 \) for practically all the wrapped monopole trajectories (while generically \( n_{\text{wrt}} \) could be equal to any integer). Moreover, the trajectories are rapidly becoming more and more static. Roughly speaking, the approximation of static trajectories is not so bad beginning with \( T = T_c \). \([21]\).

As is mentioned above, in the static limit one can consider a 3d picture. In a 3d time slice the monopole trajectories become point-like excitations which can be called instantons (in resemblance to but not in an exact correspondence with the Polyakov model \([22]\)). The density \([8]\) becomes the density of the instantons.

Within the vortices, or string \([23]\) approach the geometrical picture is similar, with the corresponding change of dimensions. At \( T > T_c \) the percolating vortices become preferably time oriented and, moreover, simply static to a reasonable approximation. In the static approximation, the 2d surfaces can be replaced by their 1d intersections with a given time-slice. The 1d defects or trajectories correspond to particles, or fields in any number of dimensions. Thus, the vortices reduce to a 3d field. While the time-direction dependence becomes trivial, the percolation in the three spatial dimensions persists. In field theoretic language this means that the corresponding 3d scalar field, \( \Sigma_M \) has a non-vanishing vacuum expectation value:

\[ \langle \Sigma_M \rangle \neq 0, \]

for further details see \([9]\).

To summarize, the 3d physics sets in for non-perturbative effects quite early, at temperatures, just above \( T_c \) the reason seems to be understood rather within dual models than within a field theoretic formulation. In the region, say,

\[ T_c < T < 2T_c \]

the parameters of the 3d field theory are to be treated phenomenologically. At larger temperatures simple scaling laws like \([9]\) become valid in many cases.

### III. TEMPERATURE DEPENDENCE

Imagine for a moment that the 3d magnetic component of plasma is indeed a liquid, as is argued on the basis of the lattice data \([8,9]\). Could it be a superconducting liquid? At first sight, the answer is obviously “no”. Indeed, ordinary superfluidity is destroyed at finite, but low temperature. But now we are discussing temperatures of order \( T_c \approx 200 \text{ MeV} \). However, why does superfluidity, present at \( T = 0 \), disappear at some \( T_0 \) despite of the fact that the two motions (superfluid and dissipative) are independent? The reason \([11]\) could be called a kind of a “nonrelativistic unitarity condition”. The density of the bosons in the condensed mode \( n_0 \) diminishes with temperature,

\[ [n_0(0) - n_0(T)]/n_0(0) \sim T^2 \quad (T \ll T_0), \]

and at \( T = T_0 \) the boson condensate gets evaporated, \( n_0(T_0) = 0 \).

In the case of Yang-Mills theory and in the weak-coupling limit, \( g^2(T) \to 0 \) the total energy/pressure density starts with the Boltzmann’s factors for non-interacting gluons. One calculates then corrections in \( g^2(T) \) and as less and less uncertainty is left in the perturbative sum the phase space available for the non-perturbative component (let it be superfluid) diminishes. However, as is mentioned above, the uncertainty of the perturbative series does not go down as an inverse power of \( T \): \([24]\) but is proportional only to \( g^2(T) \), or to \( T^3 \ln(\Lambda_{QCD}/T)^{-3} \). Thus, the r.h.s. of Eq. \((11)\) characterizes the phase space of the component which is actually not controlled by temperature and is determined by the physics in the infrared even at \( T \to \infty \).

Thus, for the phase space factor associated with the magnetic component in Eq. \((2)\) we have

\[ \langle \text{phase factor} \rangle_{\text{magnetic}} \sim \left(\frac{1}{\ln T}\right)^3, \]

which implies an amusing possibility of having superfluidity even at \( T \to \infty \) provided that the 3d field theory behind the magnetic component corresponds to a superfluid \([25]\).
IV. DYNAMICS OF THE MAGNETIC COMPONENT

The dynamics of the instantons (monopoles) in 3d Yang-Mills theories has been investigated numerically in \[26\]. In the high-temperature limit this is our magnetic component as well. One assumes that monopoles can be treated as Abelian objects with a partition function of a Coulomb gas:

\[
Z = \sum_{N=0}^{\infty} \frac{\zeta^N}{N!} \prod_{a} \int d^3 x^{(a)} \exp \left( - \frac{g^2_\text{ab}}{2} \sum_{a\neq b} q_a q_b D_{ab} \right),
\]

where \( D_{ab} \equiv D(x^{(a)} - x^{(b)}) \) is the scalar particle propagator, \( q_{a,b} \) are the monopole charges in units of elementary magnetic charge \( g_m \), \( |q_a| = 1 \), and \( \zeta \) is the fugacity. The model (13) is known to induce confinement of external electric charges \[22\].

It was found \[26\] that the lattice data can be fit by the model (13). In particular, one can define the Debye screening mass \( m_D \) of the magnetic plasma described by \[26\]. However, it turns out that \( \rho_{3,\text{mon}}/m_D^3 \approx 0.03 \) where \( \rho_{3,\text{mon}} \) is the 3d monopole density. The latter observation is in contradiction with the mechanism of the Debye screening. Another weak point of the model is that it replaces the original non-Abelian action by its Abelian projection.

In view of the observation \( \rho_{\text{mon}}/m_D^3 \ll 1 \) one might be tempted to try an opposite limit and consider the system of the 3d monopoles not as a plasma but rather as a Bose-particles system with low density. Then one of the possibility is the Bose condensation and, as a result, superfluidity \[11\]. The problem is tractable provided that the interaction region is small compared to the volume occupied by a particle on average:

\[
\rho_{\text{mon}} a^3_{\text{sc}} \ll 1 ,
\]

where \( a_{\text{sc}} \) is the scattering length. Also, the interaction is to be repulsive, \( a_{\text{sc}} > 0 \). Otherwise, the slow particles would attract each other and the condensation of the original particles is impossible. In the Abelian projection, monopoles and anti-monopoles attract each other at short distances, and the Bose condensation seems to be ruled out.

However, it is more consistent to view the “monopoles” and “antimonopoles” as non-Abelian objects detected through the Abelian projection, see, e.g., \[17\] and references therein. Then their interaction at short distances should be treated phenomenologically, through the lattice studies. It was observed \[5\] that both monopole and monopole and monopole and anti-monopole repel each other at short distances. In the language of the scattering lengths:

\[
a_{\text{mon-\text{mon}}} > 0, \quad a_{\text{mon-\text{antimon}}} > 0 ,
\]

and there is no contradiction, at this level, with the idea of the Bose condensation. A reservation is that monopole and antimonopole still attract each other at “intermediate distances”, while the monopole-monopole interaction is repulsive at all distances. The attraction, however, is not strong enough to bind the particles and in this sense can be neglected.

Moreover, it turned also possible \[6\] to extract the interaction potentials. In case of the monopole-monopole interaction the estimate is:

\[
V_{\text{mon-mon}}(r) \sim \frac{1}{r} \exp(-r/\lambda) , \quad \lambda \approx 0.1 \text{ fm} .
\]

Thus, one can estimate the interaction region as

\[
V_{\text{int}} \sim (0.2 \text{ fm})^3 .
\]

Whether the interaction region (17) is large or small depends on the density of the monopoles which is also provided by the lattice measurements \[7\].

As is mentioned above at high temperatures we expect that all the temperature dependencies are trivial. Namely, after rescaling fields and distances the 3d theory does not depend on the temperature at all. In other words, all the observables should be proportional to the corresponding powers of \( g^2(T) \cdot T \). This should be true also for the interaction region (17). However, numerically there is no much evidence for that. It is more appropriate to say that we are dealing with estimates rather than with exact numbers. The density of the monopoles, on the other hand, is measured with good accuracy. In particular,

\[
\rho_{\text{mon}} < 5 \text{ fm}^3 , \quad T_c < T < 2T_c .
\]

Thus, in this temperature range the approximation \( V_{\text{int}} \cdot \rho_{\text{mon}} \ll 1 \) seems granted. For higher temperatures the issue is more subtle and we postpone a detailed discussion of the numerics. The general impression is that the density is still low enough in the appropriate units.

Now we come to the following question. Phenomenologically, we have two descriptions of the medium of 3d excitations. If we start from the 4d monopoles at \( T = 0 \) then they become a 3d gas of instantons at \( T > T_c \). On other hand, if we start from magnetic strings then they turn into a 3d magnetically charged scalar field with non-trivial vacuum expectations value \[10\]. Both description are obtained in Abelian projections and in this sense both oversimplify the actual non-Abelian picture. However, if we are looking for a possible match of the lattice phenomenology to the theory of superfluidity, the first impression is that the pictures are mutually excluding each other and only one of them has chances to be correct, if any.

The good news is that both descriptions can correspond to superfluidity and the two pictures are just dual to each other, see, e.g., \[27\]. This duality is well known in the theory of superfluidity. One starts with the Hamiltonian for heavy particles:

\[
\hat{H} = \frac{p^2}{2m} \hat{\alpha}_{\mathbf{p}}^{+} \hat{\alpha}_{\mathbf{p}} + \frac{U_0}{2V} \sum_{\mathbf{p}_1 \neq \mathbf{p}_2} \hat{\alpha}_{\mathbf{p}_1}^{+} \hat{\alpha}_{\mathbf{p}_2}^{+} \hat{\alpha}_{\mathbf{p}_1} \hat{\alpha}_{\mathbf{p}_2} ,
\]

and
where $\hat{a}, \hat{a}^\dagger$ are annihilation and production operators. One performs then the Bogolyubov transformation,

$$\hat{a}_p = u_p \hat{b}_p + v_p \hat{b}_p^\dagger, \quad \hat{a}^\dagger_p = u_p \hat{b}^\dagger_p + v_p \hat{b}_p ,$$

(19)

where $u_p, v_p$ are coefficients, to diagonalize the Hamiltonian. In terms of the new field (associated with the operators $\hat{b}, \hat{b}^\dagger$) the spectrum starts with linear, or phonon term:

$$\epsilon(p) \approx u_p ,$$

where $u$ is the speed of sound. The new field, associated with the operators $\hat{b}, \hat{b}^\dagger$ has vacuum expectation value and we can identify this field phenomenologically with the field $\Sigma_M$, see Eq. (10).

Thus, the existence of the two descriptions of the ground state, in terms of the gas of monopoles/instantons and in terms of an infinite cluster of trajectories, or vacuum expectation value (10) is in fact a strong argument in favor of the superfluidity of the magnetic component.

Note that the linear spectrum $\epsilon(p) = u \cdot p$ corresponds to a massless field in “relativistic 3d language”. This masslessness can be traced back to the magnetic $U(1)$ symmetry of Hamiltonian (13) in terms of heavy particles. As is mentioned above the actual “monopole” and “antimonopole” do not interact at short distances as a particle and antiparticle (the terminology used is rooted in the Abelian projection but the actual non-Abelian dynamics is different). As a result, the would-be massless Goldstone boson does not show up in the spectrum of the 4d theory.

Because confinement in the spatial directions is due to breaking of magnetic $U(1)$ to magnetic $\mathbb{Z}_2$ a phenomenological 3d model which seems to be more appropriate to describe the condensation of the field $\Sigma_M$ is the ‘t Hooft model [28]:

$$\mathcal{L} = \partial_\mu \phi \partial^\mu \phi^* - M^2 \phi \phi^* - \lambda (\phi \phi^*)^2 + \frac{\zeta}{2} ((\phi)^2 + (\phi^*)^2)$$

This model has the magnetic $U(1)$ broken to $\mathbb{Z}_2$ by the $\zeta$-term and incorporates 3d confinement (area law for the spatial Wilson loop). There is no Goldstone particle. The model (21) might describe condensation of the field $\Sigma_M$.

V. CONCLUSIONS

It appears that the magnetic component of the Yang-Mills plasma could provide an independent component to the fluidity of the plasma. It is most remarkable that the properties of this component, especially its viscosity, are in a way independent of the temperature. The temperature does determine the phase factor which controls the contribution of this component of plasma to the total viscosity [2] but not the partial viscosity itself. The reason is that the magnetic component is directly related to the infrared divergences known since long time in high-temperature field theory. As a result the density of, say, monopoles is not given by a Bose distribution corresponding to the overall temperature and certain mass of the monopole. Instead, it is proportional to $(g^2(T) \cdot T)^3$ at high temperatures. To adjust phenomenology to this prediction of the theory one can introduce a corresponding chemical potential [3] but this is just another demonstration that the density of the magnetic degrees of freedom is not determined by the standard high-temperature dynamics.

The properties of the magnetic component are determined by a 3d field theory. At very high temperature it should be the standard dimensionally reduced Yang-Mills theory. At intermediate temperatures, parameters of the 3d theory can be fitted phenomenologically. In particular, the magnetic component could be superfluid. To clarify whether such a possibility realizes we invoke lattice data. Because the data are obtained in Euclidean space they refer in fact to the ground state. Phenomenologically superfluidity of the magnetic component seems plausible although further data are required to make the evaluation more reliable.

The same lattice data seem to fit known Abelian models of three-dimensional confinement as far as long-distance interaction of the constituents is concerned. At short distances, the actual non-Abelian nature of the magnetic degrees of freedom is manifested and turns crucial for the self-consistency of the models. Thus, there is a perspective that the same magnetic component of the Yang-Mills plasma could explain both the 3d confinement and low viscosity of the plasma.

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