Implementation of chromomagnetic gluons in Yang-Mills thermodynamics

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(Dated: May 11, 2014)

Motivated by the recent high-precision lattice data on Yang-Mills equations of state, we propose an effective theory of SU(3) gluonic matter. The theory is constructed based on the center and scale symmetries and their dynamical breaking, so that the interplay between color-electric and color-magnetic gluons is included coherently. We suggest, that the magnetic gluon condensate changes its thermal behavior qualitatively above the critical temperature, as a consequence of matching to the dimensionally-reduced magnetic theories. We consider thermodynamics in the mean field approximation and discuss the consequences for the interaction measure.

PACS numbers: 12.38.Aw, 25.75.Nq, 11.10.Wx

1. INTRODUCTION

Non-abelian gauge theories undergo a deconfinement phase transition at finite temperature T. Their bulk asymptotic properties are successfully captured in the quasi-particle description, which can be consistently calculated in the leading-order perturbation theory [1]. However, a naive perturbative treatment in the weak coupling g is spoiled since the magnetic screening mass is dynamically generated as a ultra-soft scale g^2T [2, 3]. The magnetic sector remains non-perturbative in the high temperature phase, and consequently, the spatial string tension is non-vanishing for all temperatures [4, 5], indicating certain confining properties.

This residual interaction brings apparent deviations in equations of state (EoS) from their Stefan-Boltzmann limit at high temperature. In particular, the interaction measure I(T) is the best observable to examine dynamical breaking of scale invariance of the Yang-Mills (YM) Lagrangian. In lattice simulations of pure SU(3) YM theory the I(T)/T^2T_c^2, with the deconfinement critical temperature T_c, is nearly constant in the range T_c < T < 5T_c. This observation strongly suggests non-trivial dynamical effects [4, 5]. Several scenarios have been proposed to explain this non-perturbative nature, such as, a dimension-2 gluon condensate that generates an effective mass term of the gauge boson dynamically [6, 7], or a temperature-dependent gluon mass [8, 9], as well as matrix models through introducing an extra T^2 term [11, 12]. Beyond this temperature range the lattice data follow the results from the Hard Thermal Loop (HTL) resummed perturbation theory. Thus, a non-perturbative part in the lattice data is extracted by subtracting the HTL contribution [2]. The resultant non-perturbative part in I(T)/T^2T_c^2 is monotonically decreasing, whereas the HTL result is monotonically increasing with T. A plateau that arises in intermediate temperatures in I(T)/T^2T_c^2 can be therefore understood as resulting from the summation of those two contributions.

In this paper, we formulate an effective theory of SU(3) gluonic matter, which accounts for two dynamically different contributions, the chromomagnetic and chromoelectric gluons.

The color-magnetic sector is described by the dilaton, whose condensate reproduces the trace anomaly of the SU(3) YM theory [13]. In general, the dilaton couples also to the Polyakov loop which is the order parameter of confinement-deconfinement phase transition and belongs to the color-electric sector. Thus, the dilaton captures the thermodynamic properties around the critical point T_c, which are related both, the color-electric and color-magnetic gluons.

Thermal behavior of the magnetic gluon condensate at high temperature is found, using the three-dimensional YM theories [14–17], to be \( \langle g^2 T \rangle^4 \) [18]. We introduce this contribution to the effective dilaton potential constructed in four dimensions. We consider the EoS in this effective theory under the mean field approximation and discuss the interaction measure and its interpretation. We also associate our qualitative result with the lattice data. It turns out that the magnetic gluon condensate can be regarded as an alternative to the HTL contribution.

2. MAGNETIC CONFINEMENT

Color-electric \( \langle E \rangle \) and color-magnetic \( \langle H \rangle \) gluon condensates behave differently at finite temperature, in particular, in the deconfined phase [10]. The phase transition is essentially driven by the electric gluons. The condensate \( \langle E \rangle \) drops toward T_c and approximately vanishes above T_c. On the other hand, the magnetic condensate, stays nearly constant across the deconfinement phase transition.

Matching the spatial string tension \( \sigma_s \), calculated from the gauge-invariant correlation function of the gauge-field strengths, to that in the 3-dimensional YM theory, yields
the magnetic condensate as \[ 18 \]
\[ \langle H \rangle = c_H \left( g^2(T)T \right)^4, \] (2.1) with
\[ c_H = \frac{6}{\pi^2} \frac{\sigma_s^2}{c_m^2}. \] (2.2) The constants \( c_s \) and \( c_m \) appear in \( \sigma_s \) and in the magnetic gluon mass as
\[ \sqrt{\sigma_s(T)} = c_s g^2(T)T, \quad m_g(T) = c_m g^2(T)T. \] (2.3) For \( SU(3) \) YM theory \( c_s = 0.566 [6] \) and \( c_m = 0.491 [20] \).

The non-vanishing string tension \( \sigma_s \) may support the conjecture, that in pure YM theory hadronic states, glueballs, can survive in deconfined phase. The scalar glueballs can be introduced as the dilatons associated with the scale symmetry. Their condensate saturates the trace anomaly through the potential \[ 13 \]
\[ V_\chi = \frac{B}{\pi} \left( \frac{\chi}{\chi_0} \right)^4 \left[ \ln \left( \frac{\chi}{\chi_0} \right)^4 - 1 \right], \] (2.4) where \( B \) is the bag constant and \( \chi_0 \) is a dimensionful constant. The two parameters, \( B \) and \( \chi_0 \), are fixed to reproduce the vacuum energy density \( \mathcal{E} = \frac{1}{4}B = 0.6 \text{ GeV fm}^{-3} \) and the vacuum glueball mass \( M_c = 1.7 \text{ GeV} [21, 22] \). One finds, that \( B = (0.368 \text{ GeV})^4 \) and \( \chi_0 = 0.16 \text{ GeV} \).

In YM theories, \( Z(N_c) \) is a relevant global symmetry that characterizes the deconfinement phase transition. The Polyakov loop \( \Phi \) is an order parameter of dynamical breaking of \( Z(N_c) \) symmetry \[ 23 \]. The \( \Phi \) is introduced as a gauge invariant operator \[ \Phi = \frac{1}{N_c} \text{tr} \hat{L}, \] (2.5) with \( \hat{P} \) being the Euclidean time ordering and \( A_4 = iA_0 \), which transforms under \( Z(N_c) \) as
\[ \Phi \rightarrow z \Phi, \quad z \in Z(N_c). \] (2.6) The potential that mixes the dilaton field and the Polyakov loop should be manifestly invariant under \( Z(N_c) \) and scale transformation. For \( N_c = 3 \), its most general form is as the following \[ 20 \],
\[ V_{\text{mix}} = \chi^4 \left( G_1 \Phi \Phi + G_2 (\Phi^3 + \Phi^3) + G_3 (\Phi \Phi)^2 + \cdots \right), \] (2.7) with unknown coefficients \( G_i \).

The Polyakov loop characterizes the chromoelectric sector of gluons. The dilaton condensate contains the information on both, chromoelectric and chromomagnetic gluons.

From the lattice results on those condensates \[ 22 \], one concludes, that the electric component of the dilaton drops toward the critical point from the side of confined phase. On the other hand, in deconfined phase, the dilaton represents chromomagnetic gluodynamics.

3. EFFECTIVE MODEL

We formulate the model of gluodynamics which accounts for the interplay between chromoelectric and chromomagnetic gluons as
\[ \Omega = \Omega_g + \Omega_\Phi + V_\chi + V_{\text{mix}} + c_0. \] (3.1) The electric gluon part \( \Omega_g \) is given in the presence of a uniform gluon field \( A_0 \) as \[ 3 \],
\[ \Omega_g = 2T \int \frac{d^3p}{(2\pi)^3} \ln \left( 1 - \bar{L}_A e^{-p/T} \right), \] (3.2) with the adjoining Polyakov-loop matrix \( \bar{L}_A \), and it can be further expressed in terms of the fundamental Polyakov loop \( \Phi \) as \[ 26 \]
\[ \Omega_g = 2T \int \frac{d^3p}{(2\pi)^3} \ln \left( 1 + \sum_{n=1}^{8} C_n e^{-np/T} \right), \] (3.3) with
\[ C_1 = 1, \] (3.4)
\[ C_2 = C_6 = 2 - 27 \Phi \Phi + 27 (\Phi^3 + \Phi^3), \] (3.5)
\[ C_3 = C_5 = -2 + 27 \Phi \Phi - 81 (\Phi \Phi)^2, \] (3.6)
\[ C_4 = 2 \left[ -1 + 9 \Phi \Phi - 27 (\Phi^3 + \Phi^3) + 81 (\Phi \Phi)^2 \right]. \]

The Haar measure part is introduced as \[ 23 \]
\[ \Omega_\Phi = -a_0 T \ln \left[ 1 - 6 \Phi \Phi + 4 (\Phi^3 + \Phi^3) - 3 (\Phi \Phi)^2 \right]. \] (3.7) To formulate an effective mixing between the Polyakov loop and dilaton, we take only the first term of Eq. (2.7). Thus,
\[ V_{\text{mix}} = G \left( \frac{\chi}{\chi_0} \right)^4 \Phi \Phi. \] (3.8) In general, the coupling \( G \) can be temperature dependent, but we consider \( G \) as a constant and fix its value to reproduce the expectation value \( \langle \Phi \rangle = 0.4 \) at \( T_c \). Requiring that a first-order phase transition appears at \( T_c = 270 \text{ MeV} \) as found in the lattice results \[ 4 \], one finds that \( a_0 = (0.184 \text{ GeV})^3 \), \( c_0 = (0.244 \text{ GeV})^4 \) and \( G = (0.206 \text{ GeV})^4 \).

Under the mean field approximation, the temperature dependence of \( \langle \Phi \rangle \) and \( \chi \) are obtained from the stationary conditions for the effective potential \[ 31 \], \[ \partial \Omega / \partial \Phi = \partial \Phi / \partial \Phi = \partial \Omega / \partial \chi = 0, \] resulting in coupled gap equations \#1. The gap equation for \( \langle \Phi \rangle \) is solved numerically, whereas that for \( \chi \) can be solved analytically as
\[ \langle \chi \rangle = \chi_0 \exp \left[ -G \langle \Phi \Phi \rangle / B \right]. \] (3.9) \#1 We note that \( \langle \Phi \rangle = \langle \Phi \rangle \).
Fig. 1 shows the expectation values of $\Phi$ and $\chi$ as the solutions of the gap equations. The thermal effect of $\langle \chi \rangle$ is induced via the mixing to the Polyakov loop, which exhibits a rather weak dependence on temperature above $T/T_c \sim 1.5$.

At higher temperature, due to the dimensional reduction, the theory in four dimensions should match the three-dimensional YM theory. We postulate the following matching condition,

$$\frac{\langle \chi \rangle}{\chi_0} = \left( \frac{\langle H \rangle}{H_0} \right)^{1/4},$$

which at a certain temperature $T_{\text{match}}$, should be met with Eq. (3.7).

A constant $H_0$ in Eq. (3.8) is chosen such, that the model reproduces the 30% reduction of the non-perturbative contribution to the interaction measure $I(T)/T^4 T_c^2$ at the matching temperature to the HTL result, as observed in the lattice calculation [7]. This implies that $H_0 = (0.8 \text{ GeV})^4$ #2.

The matching temperature $T_{\text{match}}$ can be extracted from a comparison of Eqs. (3.7) and (3.8). Applying the two-loop running coupling,

$$g^{-2}(T) = 2b_0 \ln \frac{T}{\Lambda_\sigma} + b_1 \ln \left( 2 \ln \frac{T}{\Lambda_\sigma} \right),$$

$$b_0 = \frac{11}{16\pi^2}, \quad b_1 = \frac{51}{128\pi^2},$$

with $\Lambda_\sigma = 0.104 T_c$ [8], one finds, that $T_{\text{match}} \sim 2.4 T_c$ (see Fig. 2).

In the present model, the changeover in the temperature dependence of the magnetic condensate, seen in

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The matching temperature $T_{\text{match}}$ can be extracted from a comparison of Eqs. (3.7) and (3.8). Applying the two-loop running coupling,
this contribution approaches zero, already at moderate temperature, since \( \langle \Phi \rangle \to 1 \), as seen in Fig. 4. Indeed one finds, that

\[
D \to \frac{3}{T^3} \int \frac{d^3p}{(2\pi)^3} \ln \left( 1 - e^{-p/T} \right)
+ \frac{16}{T^4} \int \frac{d^3p}{(2\pi)^3} \frac{p \cdot e^{-p/T}}{1 - e^{-p/T}}
= 8\pi^2 \left( -1 \frac{1}{15} + \frac{1}{15} \right) = 0. \tag{4.3}
\]

Requirement of the non-vanishing \( D \) can be discriminated in \( I/T^2 \) since it appears as a coefficient of the quadratic, \( T^2 \) term, whereas other contributions with \( B \) and \( C \) monotonically decrease. From Eq. (4.3), one finds, that any residual contribution \( \sim T^2 \) does not show up in \( I/T^2 \). Recall, that the coefficient \( D \) in Eq. (4.2) is entirely chromoelectric, since it does not contain \( \chi \). Therefore, in order to introduce magnetic confinement effectively which yields residual interaction at high temperature, such that \( I \neq 0 \), one transmutes \( \chi \) into \( \langle H \rangle \sim (g^2(T) T^4) \) via Eq. (5.3), and applies it to Eq. (5.1). This generates a \( T^4 \) contribution, which appears from the \( \chi + V_{\text{mix}} \) part, and results in the equations of state deviating from the Stefan-Boltzmann values at high temperature. One also finds an additional contribution to the interaction measure from \( (H) \), as

\[
\delta I = -B \frac{\langle H \rangle}{H_0} + \left( 2b_0 + \frac{b_1}{b_0} \ln \left( T/\Lambda_{\sigma} \right) \right) \frac{\langle H \rangle}{g^4(T) H_0}. \tag{4.4}
\]

The first term is of order \( O(g^8) \), whereas the second is \( O(g^4) \) which is thus the leading contribution to \( I \).

The interaction measure normalized by \( T^2 T_c^2 \) is shown in Figs. 3. The \( I/T^2 T_c^2 \) is monotonically decreasing even at high temperature when no matching to the 3-dim YM is made. The magnetic contribution generates a \( T^2 \) dependence, as seen in the figure. The sum of those two contributions forms a plateau-like behavior in \( I/T^2 T_c^2 \) at moderate temperature, \( T/T_c \sim 2-4 \). This property appears due to the residual chromomagnetic interaction encoded in the dilaton, \( \chi^4 \sim H \). The resulting behavior of \( I/T^2 T_c^2 \) with temperature, seen in Fig. 3 qualitatively agrees with the latest high-precision lattice data [7]. We note that a smooth switching from the dilaton to the magnetic condensate must happen dynamically, so that thermodynamic quantities, such as the specific heat, do not experience any irregular behavior above \( T_c \).

Fig. 4 shows the interaction measure normalized by \( T^4 \). With such normalization, the impact of the magnetic contribution is not well distinguishable. Thus, in order to identify different dynamical effects contributing to the interaction measure, it is indeed more appropriate to normalize \( I(T) \) by \( T^2 T_c^2 \), as was suggested in [12].

The lattice data have confirmed also, that there is a non-vanishing \( T^2 \) contribution to the interaction measure \( I(T) \). This would correspond to a non-vanishing \( A \) in Eq. (4.1). Such a term can appear from a dynamically generated gluon mass. Imposing the scale symmetry, the expected mass term is of the following form [10],

\[
\mathcal{L}_m = \frac{1}{2} G_A^2 \left( \frac{\chi}{\chi_0} \right)^2 A_\mu A^\mu, \tag{4.5}
\]

with a certain coupling \( G_A \).

In principle, this term should be derived from the YM Lagrangian using renormalization group. By an appropriate choice of \( G_A \), such a mass term may help to better quantify lattice data. Then, if the dynamical mass is nearly constant in a certain range of temperature, the observed plateau in \( I/T^2 T_c^2 \), may emerge even without the dilaton potential [6, 12]. Consequently, this approach could be an alternative to the formulation proposed in the previous section.
5. CONCLUSIONS

We have presented an effective theory of SU(3) Yang-Mills (YM) thermodynamics implementing the major global symmetries, the center and scale symmetries, and their dynamical breaking. This naturally allows a mixing between the Polyakov loop and the dilaton field. Consequently, the magnetic confinement is effectively embedded and results in deviations of the EoS from their Stefan-Boltzmann limit at high temperature.

Also, matching to the 3-dimensional YM theory has been proposed, which leads to the gluon condensate increasing with temperature in deconfined phase. Contrary, in the conventional treatment of the dilaton condensate, there is a weak thermal behavior of the composite gluon in a wide range of temperature. This suggests, that at some temperature above $T_c$, the gluon condensate exhibits a distinct behavior on $T$. In the present model this temperature is roughly estimated as $\sim 2.4T_c$, compatible with $\sim 2T_c$ extracted from the spatial string tension [18].

We have illustrated, that the above changeover of the gluon condensate, becomes transparent in the interaction measure $I = E - 3P$ normalized by $T^2T_c^2$, rather than by $T^4$. Adopting the matching condition [8, 3], the $I/T^2T_c^2$ shows a $T^2$ raise, which is dominating at high temperature. Before reaching the matching temperature, the thermodynamics is well described by the model for the Polyakov loop and a nearly constant dilaton condensate, resulting in a monotonic decrease of $I/T^2T_c^2$ with $T$. Consequently, the sum of those two contributions, yields a plateau structure at an intermediate temperature $T/T_c \sim 2-4$. This qualitative behavior of the interaction measure is consistent with the lattice findings [7]. The role of the magnetic gluon turns out to be alternative to the HTL contribution.

The nature of the physical vacuum in YM theories can be captured by topological objects, such as magnetic monopoles and vortices [28]. In the context of hot gluon plasma, it was shown within lattice simulations [29] that the magnetic component as a topological defect affects crucially the thermodynamics in deconfined phase.

Different approaches were proposed to deal with the magnetic aspect in the topological context [30]. In our effective theory, such magnetic feature can be attributed to the relevant global symmetries embedded in the original color gauge group. It is desirable to examine to what extent this effective theory is secure in describing the non-perturbative feature at high temperature. A matching to the topological approaches could yield more reliable constraints on the Lagrangian and its parameters.

Furthermore, introducing quarks and their coupling to gluons, in the proposed theory, could provide a scheme for an effective description of QCD thermodynamics.

Acknowledgments

We acknowledge stimulating discussions with Georg Bergner, Bengt Friman, Larry McLerran and Owe Philipsen. This work has been partly supported by the BNS-7235.2010.2 (Russia) and by the Polish Science Foundation (NCN).
(2012).
[27] K. Fukushima, Phys. Lett. B 591, 277 (2004).
[28] G. ’t Hooft, hep-th/0010225.
[29] M. N. Chernodub, A. Nakamura and V. I. Zakharov, Phys. Rev. D 78, 074021 (2008). M. N. Chernodub, Y. Nakagawa, A. Nakamura, T. Saito and V. I. Zakharov, Phys. Rev. D 83, 114501 (2011).
[30] P. Giovannangeli and C. P. Korthals Altes, Nucl. Phys. B 608, 203 (2001). J. Liao and E. Shuryak, Phys. Rev. C 75, 054907 (2007). M. N. Chernodub and V. I. Zakharov, Phys. Rev. Lett. 98, 082002 (2007).