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A novel massless mode in hot QCD

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I present recent results on a novel massless mode generated by the (chromo)magnetic scale g^2T in hot QCD which induces genuine non-Abelian effects such as positivity violation.

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1. Introduction

Conventional thermal perturbation theory breaks down at the (chromo)magnetic scale g^2T due to the Linde problem [1, 2]. The nonperturbative nature of the magnetic scale is intimately related to the confining property of the dimensionally reduced Yang-Mills theory at high temperature. This suggests that a confinement mechanism should be incorporated within perturbative expansions even when dealing with the deconfined quark-gluon plasma (QGP) phase. Color confinement is deeply related to *positivity violation* of the spectral function: if the spectral function of a particle is not positive semi-definite, no Källén-Lehmann representation exists, it is then not part of the physical spectrum and thus confined (see Refs. [3, 4] for reviews). Conventional thermal field approaches to hot QCD are based on massive quasiparticles which only generate short-range correlations (see Refs. [5, 6, 7] for reviews). In order to describe a strongly coupled QGP as produced in heavy-ion collision experiments, long-range correlations, whose carriers are light and/or massless modes, are a crucial ingredient. In the following, I briefly discuss the first study on massless modes in hot QCD using confining gluons recently reported in Ref. [8] which shows genuine non-Abelian features such as positivity violation.

A formalism to tackle the issue is the Gribov-Zwanziger (GZ) action, which is well-known from the study of color confinement [9, 10]. It regulates the IR behavior of QCD by fixing the residual gauge transformations, i.e., Gribov copies, that remain after applying the Faddeev-Popov procedure. The GZ action is renormalizable, and it thus provides a systematic framework for perturbative calculations (i.e., $g \ll 1$) incorporating confinement effects. The gluon propagator in general covariant gauge reads

$$D^{\mu\nu}(P) = \left[\delta^{\mu\nu} - (1-\xi)\frac{P^\mu P^\nu}{P^2}\right]\frac{P^2}{P^4+\gamma_G^4}\,, \tag{1.1}$$

where ξ is the gauge parameter. The Gribov parameter γ_G is solved self-consistently from a gap equation that is defined to infinite loop orders. The GZ gluon propagator is IR suppressed, manifesting confinement effects, and it is a significant improvement over the one from the Faddeev-Popov quantization which forms the basis for conventional perturbative calculations. The gap equation at one-loop order can be solved analytically at asymptotically high T and gives [11, 12]

$$\gamma_G = \frac{D-1}{D}\frac{N_c}{4\sqrt{2}\pi}g^2T\,, \tag{1.2}$$

where D is the space-time dimensions and N_c is the number of colors. Eq. (1.2) provides a fundamental IR cutoff at the magnetic scale for the finite-T GZ action.

2. Results and discussions

An important measure for the collective behavior of a QGP is the self-energy of quarks and gluons, from which thermal masses, dispersion relations, and spectral functions of collective excitations are derived. The Euclidean one-loop quark self-energy reads

$$\Sigma(P) = (ig)^2 C_F \sum\!\!\!\!\!\!\!\int_{\{K\}} \gamma^\mu S(K)\gamma^\nu D^{\mu\nu}(P-K)\,, \tag{2.1}$$

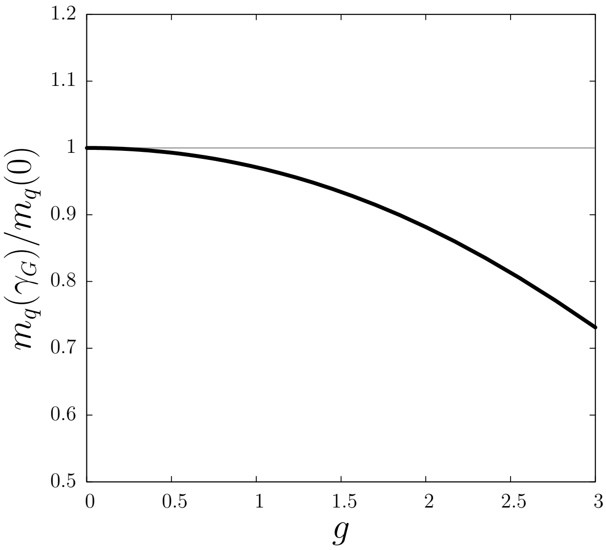

Figure 1: The quark thermal mass $m_q(\gamma_G)$ scaled by the perturbative value $m_q(0)$.

where $S(P) = 1/\not{P}$ is the quark propagator, and $D^{\mu\nu}(P)$ is the gluon propagator which is taken from Eq. (1.1). It is worth noting that there have been similar studies for the quark self-energy with nonperturbative gluons at finite density [13, 14] and in strong magnetic fields [15].

Following the systematics of the hard-thermal-loop (HTL) effective theory [16], the gauge-invariant contribution to Eq. (2.1) reads [8]

$$\Sigma(P) \simeq -(ig)^2 C_F \sum_{\pm} \int_0^\infty \frac{\mathrm{d}k}{2\pi^2} k^2 \int \frac{\mathrm{d}\Omega}{4\pi} \frac{\tilde{n}_\pm(k,\gamma_G)}{4E^0_\pm} \left[\frac{i\gamma_0 + \hat{k}\cdot\gamma}{iP_0 + k - E^0_\pm + \frac{p\cdot k}{E^0_\pm}} + \frac{i\gamma_0 - \hat{k}\cdot\gamma}{iP_0 - k + E^0_\pm - \frac{p\cdot k}{E^0_\pm}} \right], \tag{2.2}$$

where $\hat{k} = k/k$ with $k = |k|$, $E^0_\pm = \sqrt{k^2 \pm i\gamma_G^2}$, $\tilde{n}_\pm(k,\gamma_G) \equiv n_B(\sqrt{k^2 \pm i\gamma_G^2}) + n_F(k)$ with n_B and n_F the Bose-Einstein and Fermi-Dirac distributions, and $\int \mathrm{d}\Omega = \int_0^{2\pi} \mathrm{d}\phi \int_0^\pi \mathrm{d}\cos\theta$.

The quark thermal mass incorporating effects from the magnetic scale reads

$$m_q^2(\gamma_G) = \frac{g^2 C_F}{4\pi^2} \sum_\pm \int_0^\infty dk \frac{k^2 \tilde{n}_\pm(k,\gamma_G)}{E^0_\pm}, \tag{2.3}$$

which reduces to the conventional HTL one, $m_q^2(0) = C_F g^2 T^2/8$, for $\gamma_G = 0$. The scaled quark thermal mass $m_q(\gamma_G)/m_q(0)$ is shown in Fig. 1. It is clear from the figure that m_q receives negative contributions from γ_G, which is a manifestation of anti-screening effects generated by the magnetic scale. Although the effect is modest in the studied range of couplings, this is a profound signal of the build-up of long-range correlations in the system and similar anti-screening effects have been observed on the lattice for the Debye screening mass [17].

The dispersion relation is obtained by analytically continuing the self-energy Eq. (2.2) to Minkowski space and then solving the poles in the corresponding quark propagator $iS^{-1}(P) = \not{P} - \Sigma(P) = 0$. The resulting dispersion relations and residues of the poles are displayed in the upper and lower panels of Fig. 2.

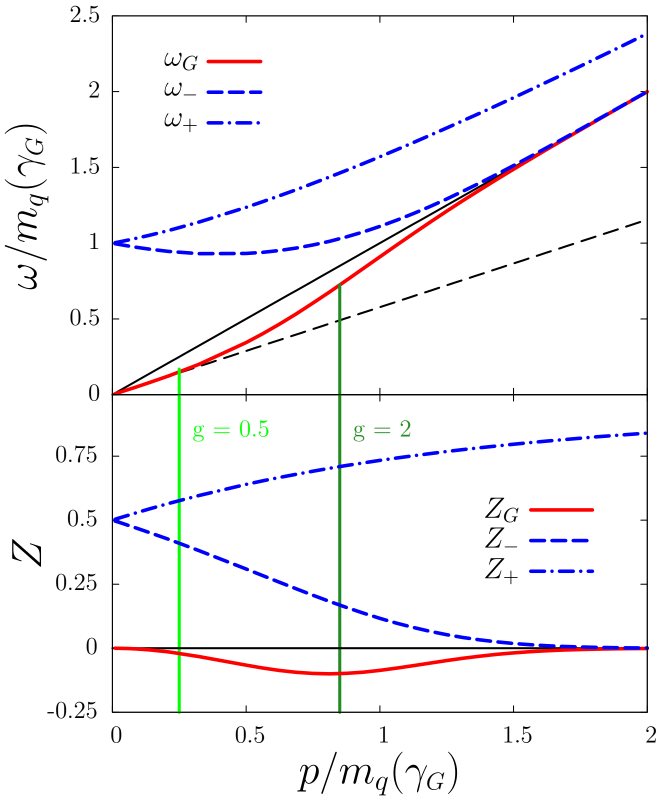

Figure 2: Dispersion relations (upper panel) and the corresponding residues (lower panel) for the particle (ω_+), plasmino (ω_-) and Gribov (ω_G) poles.

In contrast to the conventional HTL case, there are three poles in the propagator. Firstly, the screened quasi-particle excitations are recovered,

$$\omega = \omega_+(p;\gamma_G)\,, \qquad \omega = \omega_-(p;\gamma_G)\,, \tag{2.4}$$

the so-called particle ω_+ and plasmino ω_- modes, with $\omega_\pm(0;\gamma_G) = m_q(\gamma_G)$ as expected. Both $\omega_\pm/m_q(\gamma_G)$ and $Z_\pm$ are g-independent, and this has been verified explicitly up to $g \sim 2$ in Ref. [8]. This property is exactly the same as in the conventional HTL effective theory, and it is thus a non-trivial consistency check of the setup.

In addition to the massive modes, there exists a novel excitation named *Gribov pole* as in Ref. [8],

$$\omega = \omega_G(p;\gamma_G)\,. \tag{2.5}$$

It describes *massless* fermionic excitations in the plasma with dispersion relation $\omega = v_s p$ at small momenta, with $v_s \approx 1/\sqrt{3}$ (speed of sound) independent of g for the studied range. The Gribov mode "grows" in the (ω, p)-plane while the magnetic scale is increasing (through increasing g), and this effectively introduces a new *magnetic scaling* behavior to the non-Abelian plasma. The vertical lines in Fig. 2 schematically demonstrate how the Gribov mode grows: at small coupling, e.g., $g = 0.5$, the mode terminates at rather small momentum; as the coupling increases, to e.g., $g = 2$, the permitted momentum range inceases accordingly. At larger momenta than the permitted ones for each coupling, we are hitting branch cuts and Landau damping takes place as a consequence. The Gribov pole goes along with a residue $Z_G(p) < 0$ which directly implies *positivity violation* of the corresponding spectral functions in the region of space-like momenta. These novel features are direct manifestations of long-range confinement effects surviving at finite T in the non-Abelian plasma. The results reflect common features of Gribov-like approaches [9, 10, 18], though the calculation was done via the GZ action.

The uncovering of the massless Gribov mode has been an exciting attempt in exploring the significance of the magnetic scale to a non-Abelian plasma. It sheds new light on the active degrees of freedom released in course of a heavy-ion collision through which a strongly coupled QGP might emerge. It would be extremely tempting and challenging to explore the phenomenological significance of this new mode in interpreting experimental data.

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