Screening masses in a gluon plasma : a non-perturbative appraisal

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One of the basic signatures of the deconfined plasma phase of quarks and gluons is the screening of static chromoelectric fields. In the deconfined phase, the the potential between quark and antiquark pairs are exponentially screened [1] which allows forensic study of quark-gluon plasma (QGP) created in the laboratory through the measurement of the relative abundance of heavy quark bound states such as J/ψ for example [2].

In perturbation theory, the lowest order electric screening mass is obtained as $m_D \sim gT$ (the strong coupling constant is $\alpha_s = g^2/4\pi$ and T is the temperature). Magnetic screening mass, on the other hand, is expected to be generated ($\sim gm_D$) nonperturbatively in the static sector [3]. Moreover, the perturbative methods could only be reliable for temperatures far above the critical temperature ($T > T_c$) when the coupling is small ($g \to 0$).

It is widely accepted that the nonperturbative dynamics of QCD is signaled by the emergence of power corrections in physical observable. These nonperturbative corrections are introduced via non-vanishing vacuum expectation values of local quark and gluonic operators such as the quark condensate $\langle\bar{\psi}\psi\rangle$ and the gluon condensate $\langle\mathcal{G}^a_{\mu\nu}\mathcal{G}^{a\mu\nu}\rangle$. This approach of the operator product expansion (OPE) has met noticeable success in QCD sum rule calculations at zero temperature [4] and the calculation of the N-point functions in QCD at zero temperature [5]. The aim of this work is to determine the screening masses in a gluonic heatbath using operator product expansion.

The in-medium gluon self-energy for an $O(3)$ invariant gauge fixing condition can be written as

$$\pi_{\mu\nu}(\omega, k) = \pi_l \mathcal{P}^l_{\mu\nu} + \pi_t \mathcal{P}^t_{\mu\nu} + \pi_m \mathcal{M}_{\mu\nu} + \tilde{\pi}\mathcal{L}_{\mu\nu}\,. \tag{1}$$

The Lorentz invariant single particle energy and momentum are, respectively, given as $\omega = u\cdot K$ and $k = \sqrt{(u\cdot K)^2 - K^2}$ where u^μ is the four velocity of the heat bath and $K = \left(k_0, \vec{k}\right)$. The projection operators chosen here are similar to that in [6] with $\widetilde{K_\mu} = K_\mu - \omega u_\mu$ and $\bar{u}^\mu = u^\mu - (\omega/K^2)K^\mu$. The scalar functions in (1) are extracted as

$$\pi_l = \mathcal{P}^{\mu\nu}_l \pi_{\mu\nu}\,,\ \pi_t = \frac{1}{2}\mathcal{P}^{\mu\nu}_t \pi_{\mu\nu}\,,$$
$$\pi_m = -\mathcal{M}^{\mu\nu}\pi_{\mu\nu}\,,\ \widetilde{\pi} = \mathcal{L}^{\mu\nu}\pi_{\mu\nu}\,.$$

and these functions receive both perturbative and nonperturbative contributions.

In OPE, the nonperturbative corrections to the polarization tensor are calculated by writing down the full Feynman diagrams and subtracting the equivalent perturbative ones. The soft loop momenta are expanded in powers of external momenta and moments are identified with appropriate condensates. This is different from the Hard Thermal Loop approximation of pQCD where the polarization operator is saturated by the hard loop momenta ($\sim T$). The nonperturbative contribution to various scalar components of the gluon-self energy in the static limit $k_0 \to 0$ can be obtained as

$$\pi_l^{np}(0,k) = -\frac{a}{k^2}\,,$$
$$\pi_t^{np}(0,k) = -\frac{b}{k^2}$$
$$\pi_m^{np}(0,k) = 0, \qquad \tilde{\pi}^{np}(0,k) \neq 0\,, \quad (3)$$

where,

$$a = \frac{4\pi^2 N_c}{F_A}\left[\frac{8}{3}\frac{\alpha_s}{\pi}\left\langle \mathcal{E}^2\right\rangle_T + \frac{8}{30}\frac{\alpha_s}{\pi}\left\langle \mathcal{B}^2\right\rangle_T\right], \quad (4a)$$

$$b = \frac{4\pi^2 N_c}{F_A}\left[W_E\frac{\alpha_s}{\pi}\left\langle \mathcal{E}^2\right\rangle_T - W_B\frac{\alpha_s}{\pi}\left\langle \mathcal{B}^2\right\rangle_T\right]. \quad (4b)$$

Here $F_A = N_c^2 - 1$, $W_E = \left(2 + \frac{\xi}{3}\right)$, $W_B = \frac{1}{15}(38 + 9\xi)$ and $W_M = (2 + \xi)$. N_c is the number of color. The electric and magnetic screening masses are obtained from

$$m_e^2 = \pi_l^{pert}\left(0, -m_D^2\right) + \pi_l^{np}\left(0, -m_D^2\right) \quad (5)$$
$$m_m^2 = \pi_t^{pert}\left(0, -m_m^2\right) + \pi_t^{np}\left(0, -m_m^2\right) \quad (6)$$

To the perturbative order α_s, $\pi_l^{pert}(0,k) = (m_e^{pert})^2 = 4\pi\alpha_s T^2$, whereas $\pi_t^{pert}(0,p) = 0$. Solving (5) and (6), we obtain the values of the screening masses as

$$m_m = b^{\frac{1}{4}}\,,$$
$$m_e = \left[\frac{1}{2}\left\{(m_e^{pert})^2 + \sqrt{(m_e^{pert})^4 + 4a}\right\}\right]^{\frac{1}{2}} \quad (7)$$

For numerical evaluations of nonperturbative part in screening masses we use electric and magnetic condensates measured on lattice for pure $SU(3)$ gauge theory [7]. The perturbative part is evaluated using the two loop running coupling constant.

The electric and magnetic screening masses so obtained are delineated in Fig. 1. The complete electric screening mass, including perturbative and nonperturbative contributions, falls short of lattice data but is still rather close to it. On the other hand, the magnetic screening is purely nonperturbative in nature and agrees relatively well with lattice

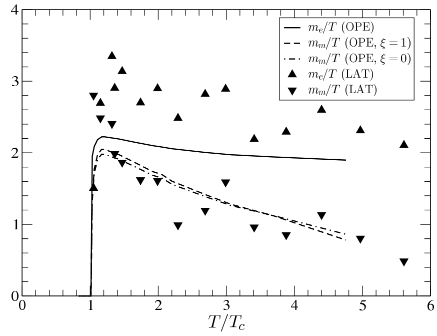

FIG. 1: Temperature variation of electric and magnetic screening masses. The present investigation is represented by OPE with two different gauge fixing parameter ξ whereas LQCD data are represented by LAT [8].

data. The magnetic mass is dependent upon the gauge fixing parameter and we have chosen Landau ($\xi = 0$) and Feynman ($\xi = 1$) gauge. The weak gauge dependence found here is similar to that of Ref. [8] for a similar choice of gauge fixing.

Our results may be useful input to calculate various thermodynamic quantities, spectral properties and for the phenomenology of jet quenching, quarkonium suppression in hot QCD matter produced in relativistic heavy-ion collision experiments.

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