Gluon quasi-particle in chirally imbalanced medium

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There is a multitude of degenerate vacuum states in Quantum Chromodynamics (QCD), each characterized by an integer-valued winding number and separated by potential energy barriers. At lower temperatures, these vacuum states can be explored through non-trivial topological gauge field configurations known as instantons. However, at high temperatures, such as in the quark-gluon plasma (QGP) phase, there is an expectation of abundant production of another type of gluon configuration called sphalerons. These topologically non-trivial gauge field configurations have the ability to change the helicity of quarks during interactions. This leads to the localized breaking of parity (P) and charge-parity (CP) symmetries through the axial anomaly of QCD, creating an asymmetry between left and right-handed quarks. This chirality imbalance is quantified using a chiral chemical potential (CCP) denoted as μ_5, which essentially represents the difference between the numbers of right and left-handed quarks. Chiral systems hold significant relevance across various fields, including particle physics, nuclear physics, condensed matter physics, and cosmology. A notable example is the chiral magnetic effect (CME), which involves generation of a parity-violating current within a plasma when exposed to a magnetic field.

When elementary particles move through a medium, their properties change due to interactions, a process known as becoming 'dressed.' This includes acquiring effective masses different from those in vacuum. Additionally, the propagation can involve collective modes or quasi-particles, which often have dispersion relations and behaviors distinct from those in vacuum, leading to novel dispersive modes. In this work we are interested in the collective oscillations of gluons in the presence of a chirally asymmetric medium.

The general structure of gluon self-energy in a chirally imbalanced medium is given by [1]

$$\Pi^{\mu\nu} = \Pi_T R^{\mu\nu} + \Pi_L Q^{\mu\nu} + \Pi_P P^{\mu\nu} \tag{1}$$

where Π_P is purely finite CCP effect. Expressions for the basis tensors can be found in [1]. Now employing the Dyson-Schwinger equation, one can write the general structure of effective gluon propagator as

$$D^{\mu\nu} = \frac{A^{\mu\nu}}{K^2 + \Pi_T + \Pi_P} + \frac{B^{\mu\nu}}{K^2 + \Pi_T - \Pi_P} + \frac{Q^{\mu\nu}}{K^2 + \Pi_L} + \frac{\xi}{K^2}\frac{K^\mu K^\nu}{K^2} \tag{2}$$

where K in the 4-momentum of the external gluon and $A^{\mu\nu} = (R+P)^{\mu\nu}/2, B^{\mu\nu} = (R-P)^{\mu\nu}/2$. The collective modes of gluon in a chirally asymmetric matter can be found from the poles of the complete propagator given by Eq. (2). Thus the dispersion relations are

$$K^2 + \Pi_T \pm \Pi_P = 0 \tag{3}$$

$$K^2 + \Pi_L = 0 \tag{4}$$

To achieve this it is necessary to evaluate all the structure factors $\Pi_{T,L,P}$ using the orthogonality properties of the basis tensors which yields [1]

$$\Pi_T = \frac{1}{2} R^{\mu\nu}\Pi_{\mu\nu},\ \Pi_L = Q^{\mu\nu}\Pi_{\mu\nu}, \quad \Pi_P = -\frac{1}{2} P^{\mu\nu}\Pi_{\mu\nu}\ . \tag{5}$$

Thus it boils down to the calculation of $\Pi^{\mu\nu}$. Here we will take the real time approach of

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thermal field theory in which the thermal propagator assumes 2×2 matrix form. However, the knowledge of only 11-component of this matrix is sufficient for our purpose, which is given by

$$\Pi_{11}^{\mu\nu} = -ig^2 \int \frac{d^4P}{(2\pi)^4} \mathrm{Tr}\left[S_{11}(P)\gamma^\mu S_{11}(K')\gamma^\nu\right] \tag{6}$$

where $K' = K + P$. Here S_{11} is the 11-component of the fermion propagator at finite temperature and CCP [2]. Now using Eqs. (5) and (6) one can calculate all the non-trivial structure factors and under Hard thermal loop (HTL) approximation they are given by

$$\begin{aligned}
\Pi_T &= -\frac{m_D^2}{2}\frac{k_0^2}{k^2}\left[1+\frac{1}{2}\left(\frac{k}{k_0}-\frac{k_0}{k}\right)\log\frac{k_0+k}{k_0-k}\right] \\
\Pi_L &= m_D^2\frac{k_0^2-k^2}{k^2}\left[1-\frac{k_0}{2k}\log\frac{k_0+k}{k_0-k}\right] \\
\Pi_P &= -\frac{g^2\mu_5}{2\pi^2}\frac{k_0^2-k^2}{k^2}\left[1-\frac{k_0}{2k}\log\frac{k_0+k}{k_0-k}\right]
\end{aligned} \tag{7}$$

where $m_D^2 = g^2\left[T^2/3 + (\mu_R^2+\mu_L^2)/2\pi^2\right]$. One can use these expressions for $\Pi_{T,L,P}$ to calculate different dispersive modes using Eqs. (3)-(4).

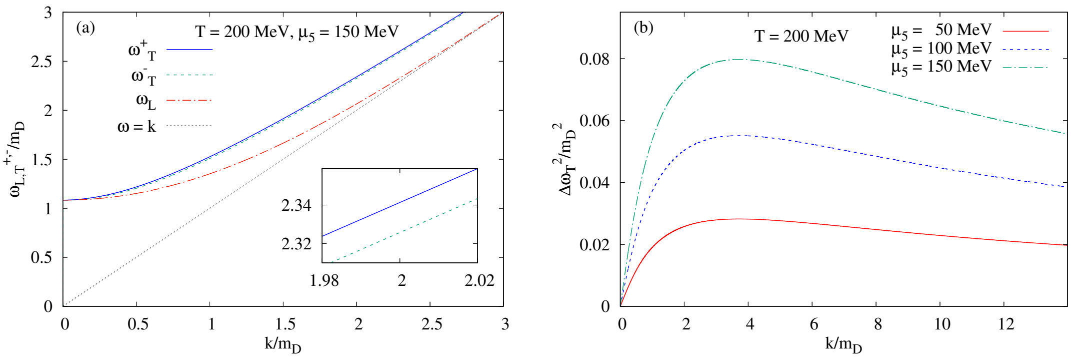

Figure 1: (a) Dispersion of $\omega_{T,\pm}$ and ω_L modes and (b) difference between two transverse modes are plotted as function of k. All results are properly scaled by debye mass m_D to make them dimensionless.

In this work we will only study real solutions of the dispersion Eqs. (3)-(4) which are shown in Fig. 1 (a). In addition to these real solutions, there are purely imaginary poles leading to the chiral plasma instability. Note that in absence of chiral imbalance i.e. $\mu_5 = 0$, the collective mode associated with Eq. (3) is called the transverse mode which is doubly degenerate [3]. However, finite values of μ_5 lifts this degeneracy and two different transverse stable collective modes are observed denoted as ω_T^+ and ω_T^- (see the inset plot where zoomed version of these two modes are shown). The solution of Eq. (4) gives the longitudinal mode of propagation with energy known as plasmon, denoted by ω_L. It is a long wavelength mode that arises solely due to the presence of the thermal medium and reduces to free dispersion faster than the transverse one by decoupling from the plasma. Fig. 1 (b) illustrates the difference between two non-degenerate transverse modes as a function of k for different values of μ_5. Initially, the two modes are degenerate at $k = 0$. As k increases, the splitting between the modes grows, reaching a maximum around $k \approx 2m_D$, after which the splitting begins to decrease.

References

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