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What can we learn from electromagnetic plasmas about the quark–gluon plasma?

Markus H Thoma

Max-Planck-Institut für extraterrestrische Physik, PO Box 1312, 85741 Garching, Germany

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Abstract

Ultra-relativistic electromagnetic plasmas can be used for improving our understanding of the quark–gluon plasma. In the weakly coupled regime, both plasmas can be described by transport theoretical and quantum field theoretical methods leading to similar results for the plasma properties (dielectric tensor, dispersion relations, plasma frequency, Debye screening, transport coefficients, damping and particle production rates). In particular, future experiments with ultra-relativistic electron–positron plasmas in ultra-strong laser fields might open the possibility of testing these predictions, e.g. the existence of a new fermionic plasma wave (plasmino). In the strongly coupled regime, electromagnetic plasmas such as complex plasmas can be used as models or at least analogies for the quark–gluon plasma possibly produced in relativistic heavy-ion experiments. For example, pair correlation functions can be used to investigate the equation of state and cross section enhancement for parton scattering can be explained.

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1. Introduction

At very high temperatures ($T > 150$ MeV) and densities ($\rho > 1\text{ GeV fm}^{-3}$), it was predicted that QCD implies that there should be a phase transition from nuclear or hadronic matter to a system of deconfined quarks and gluons [1]. The early Universe should have been in this state for the first few microseconds [1]. Such a phase transition could also occur in the interior of neutron stars [2] or in relativistic nucleus–nucleus collisions leading to the so-called quark–gluon plasma (QGP) [3]. The phase diagram is sketched in figure 1. The nature of this phase transition is similar to the Mott transition from an insulator to a metal [4], where the insulator phase at low pressure corresponds to the confined phase (nuclear matter) and the metallic phase at high pressure, e.g. metallic hydrogen [5] predicted to exist in the interior of Jupiter [6], to the deconfined phase (QGP).

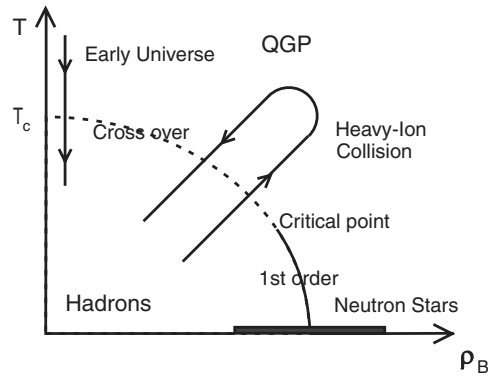


Figure 1. Sketch of the QCD phase diagram.

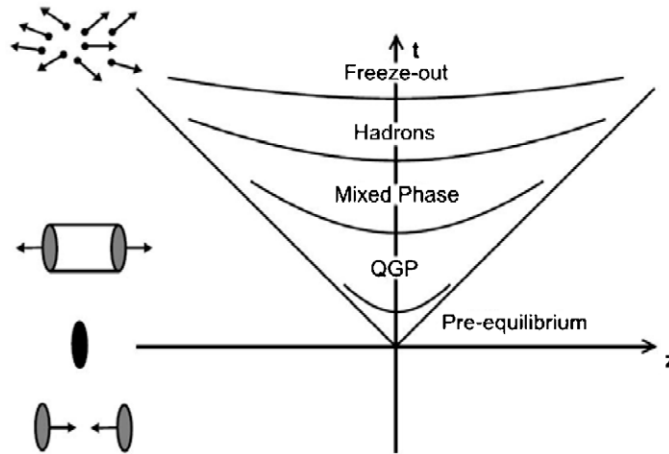


Figure 2. Spacetime evolution of the fireball in an ultra-relativistic heavy-ion collision.

Estimates show that a QGP droplet containing a few thousand quarks and gluons (partons) can be formed in the fireball of a heavy-ion collision which exists for a few $\text{fm} c^{-1}$. The spacetime evolution of the fireball looks like this (see figure 2) [7]: after the collision of the two nuclei, in which the nucleons dissolve into quarks and gluons, a pre-equilibrium phase is present for about $1\text{fm} c^{-1}$ which approaches equilibrium by secondary parton interactions. If the temperature and energy density are above the critical values, the fireball will be in the QGP phase, which should be in thermal but may be not in chemical equilibrium [8]. Due to the expansion, the fireball will cool and the transition temperature will be reached after a few $\text{fm} c^{-1}$ followed by a mixed phase in the case of a first-order transition and a hadronic phase. Finally, the hadrons will cease to interact if the system becomes dilute (freeze-out).

Since the QGP cannot be observed directly, the big question is how to detect it. Only by comparing theoretical predictions of signatures for the QGP formation with experimental data, i.e. by circumstantial evidence, can this goal be achieved. Hence, a quantitative and detailed theoretical description of the QGP is required. Basically, there are three different methods: the first one, perturbative QCD, is valid only for large parton momenta or extreme

high temperatures for which the interaction between the partons becomes weak due to a property of QCD called asymptotic freedom. Perturbative QCD allows the calculation of static properties, e.g. equation of state, as well as dynamical quantities, e.g. parton scattering cross sections. The second method, lattice QCD, is a truly non-perturbative method based on numerical simulations for solving the QCD equations on a discrete four-dimensional spacetime lattice. Unfortunately, this technique can only be applied to static quantities whereas most signatures of the QGP are dynamical like particle production rates. The third method, which I would like to discuss here in particular, is to apply and extend techniques from classical electromagnetic plasmas to the QGP. For example, transport theory (Boltzmann equation) or molecular dynamical simulations are widely used to describe properties of plasmas. Also, analogies with ideas and analytic models in plasma physics can be useful. I will explain this in the following sections following partly the review article [9].

It should be noted, however, that there are important differences between the QGP and usual ion–electron plasmas. First of all, the QGP is a relativistic system. Therefore a comparison with a relativistic electron–positron plasma, as is found in supernova explosions, is more appropriate. Second, the interactions in the QGP are based on QCD, i.e. color charges instead of electric charges have to be considered. In particular, the non-Abelian character of the interaction between quarks and gluons has a strong influence on the properties of the QGP especially at temperatures close to the deconfinement transition. Hence, the comparison with electromagnetic plasmas allows only qualitative insights for the properties of the QGP in ultra-relativistic heavy-ion collisions.

First I will start with the weakly coupled phase of the QGP, in which perturbative QCD can be used. As we will see, the methods and results are very similar to a hot QED plasma (electron–positron plasma) which can be investigated in the laboratory using ultra-strong lasers in the near future and can therefore serve also as a test model for the QGP. Afterward, I will discuss the strongly interacting phase of the QGP by comparing it with strongly coupled electromagnetic plasmas as discussed in [10–12].

2. The weakly coupled quark–gluon plasma

Interactions between partons in the QGP lead to collective phenomena, such as Debye screening and plasma waves, or transport properties, such as viscosity. At temperatures far above the critical temperature T_c for the deconfinement phase transition, the effective temperature-dependent strong coupling constant, $\alpha_s = g^2/4\pi$, becomes small. At temperatures which can be reached in heavy-ion experiments ($T < 4T_c$), $\alpha_s = 0.3$ – 0.5 is not really small rendering the applicability of perturbation theory, which is an expansion in the coupling constant, questionable. Perturbation theory in quantum field theory is most conveniently done by using Feynman diagrams which can be via Feynman rules directly translated into scattering amplitudes from which physical measurable quantities such as cross sections, damping and production rates, and lifetimes follow. If the interactions take place in the presence of a heat bath such as the QGP background, one has to consider QCD at finite temperature. For this purpose the Feynman rules have to be generalized to finite temperatures (and chemical potential), which can be done in either the imaginary [13] or real [14] time formalism. For an application of this method to the QGP, see, for example, [15].

Alternatively, transport theory can be used. For example, the dielectric tensor of the QGP can be derived by combining the Vlasov and Maxwell equations. In an isotropic QED plasma, in which there are two independent components, a longitudinal and a transverse, one finds in

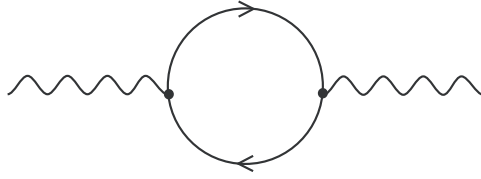


Figure 3. Polarization tensor in lowest order perturbation theory.

this way [16]

$$\begin{aligned} \epsilon_L(\omega, k) &= 1 + \frac{3m_\gamma^2}{k^2} \left[1 - \frac{\omega}{2k} \ln \frac{\omega + k}{\omega - k} \right], \\ \epsilon_T(\omega, k) &= 1 - \frac{3m_\gamma^2}{2k^2} \left[1 - \left(1 - \frac{k^2}{\omega^2} \right) \frac{\omega}{2k} \ln \frac{\omega + k}{\omega - k} \right], \end{aligned} \quad (1)$$

where ω is the frequency, $k = |\mathbf{k}|$ is the momentum and $m_\gamma = eT/3$ is the plasma frequency. In the case of the QGP, this result also holds if one replaces m_γ by $m_g = \sqrt{(1 + n_F/6)/3} gT$, where n_F is the number of quark flavors in the QGP. The Debye screening length following from the static limit ($\omega = 0$) of the longitudinal dielectric function is given by $\lambda_D = 1/(\sqrt{3}m_\gamma)$. For $\omega^2 < k^2$, the dielectric functions become negative corresponding to Landau damping.

The same result can be obtained by calculating the polarization tensor $\Pi_{\mu\nu}$ of figure 3 in the high-temperature approximation [17, 18] and using the relation [19]

$$\epsilon_L(\omega, k) = 1 - \frac{\Pi_L(\omega, k)}{k^2}, \quad \epsilon_T(\omega, k) = 1 - \frac{\Pi_T(\omega, k)}{\omega^2}, \quad (2)$$

where $\Pi_L = \Pi_{00}$ and $\Pi_T = \sum_{ij} (\delta_{ij} - k_i k_j / k^2) \Pi_{ij} / 2$ with i and j representing the space indices.

Using the Maxwell equations the dispersion relations $\omega_{L,T}(k)$ of plasma waves, describing the propagation of electromagnetic, i.e. photons (or chromoelectromagnetic, i.e. gluons, in the case of a QGP), waves in the plasma, can be found from

$$\epsilon_L(\omega, k) = 0, \quad \epsilon_T(\omega, k) = \frac{k^2}{\omega^2}. \quad (3)$$

The dispersion relations are shown in figure 4, where the branch ω_L is called a plasmon, representing a longitudinal electromagnetic wave in the medium which does not exist in vacuum. In a relativistic plasma the transverse waves associated with the magnetic interaction are as important as the longitudinal ones (plasmons), whereas in non-relativistic plasmas only the plasmons are of significance. If one wants to go beyond the high-temperature approximation, one cannot use classical transport theory but has to consider higher order diagrams for the polarization tensor.

Further interesting quantities can be derived from the dielectric tensor. For example, the wake potential of a moving charge Q with velocity \mathbf{v} , such as a heavy quark propagating through the QGP with a large transverse momentum from initial hard collisions, is given by [20]

$$\phi(\mathbf{r}, t, \mathbf{v}) = \frac{Q}{2\pi^2} \int d^3k \frac{e^{-i\mathbf{k} \cdot (\mathbf{r} - \mathbf{v}t)}}{k^2 \epsilon_L(\omega = \mathbf{v} \cdot \mathbf{k}, k)}. \quad (4)$$

Wakes created by fast quarks in a QGP lead to an attraction and the possible formation of diquarks [21]. Wakes and Mach cones have also been observed in complex or dusty plasmas

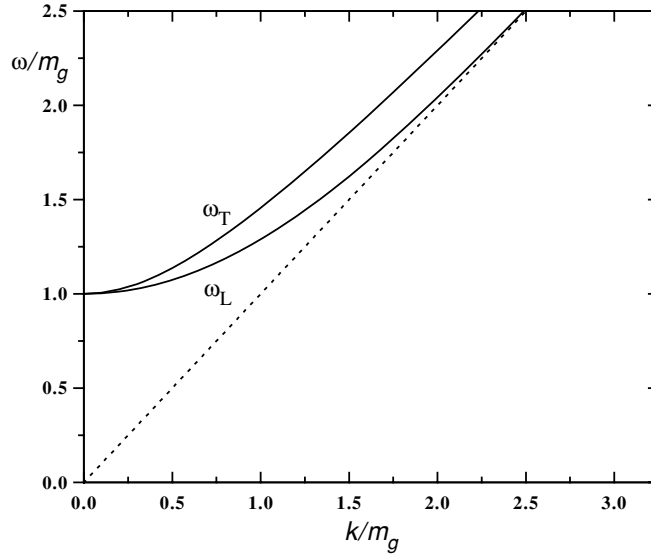


Figure 4. Dispersion relations of photons in a QED plasma or gluons in a QGP.

[22]. A closely related quantity is the energy loss of an energetic quark or gluon in the QGP. Partons with a large transverse momentum p_T are created from initial hard collisions between the partons in the nuclei. These quarks and gluons propagate through the fireball losing energy. The energy loss coming from the back reaction of the induced electric field of a fast moving charge in a plasma can be expressed by the dielectric functions of the plasma [23]. In the case of a high energy quark with velocity v in the QGP, this energy loss is given by [24]

$$\frac{dE}{dx} = -\frac{4\alpha_s}{3\pi v^2} \int \frac{dk}{k} \int_{-vk}^{vk} d\omega \omega \left[\text{Im} \frac{1}{\epsilon_L(\omega, k)} + (v^2 k^2 - \omega^2) \text{Im} \frac{1}{\omega^2 \epsilon_T(\omega, k) - k^2} \right]. \quad (5)$$

In addition, there are contributions to the energy loss from elastic parton scattering and from bremsstrahlung (radiative energy loss) [25]. The combination of the energy loss by the induced electric field, corresponding to long-range interactions within the plasma, together with individual elastic scatterings is called collisional energy loss. A consistent perturbative treatment of the collisional energy loss based on the hard thermal loop resummation technique [26] was presented in [27]. Later on, it was argued that the radiative energy loss dominates for relativistic partons in the QGP (see e.g. [25]). However, it could be shown [28, 29] that for the quenching of high- p_T hadron spectra resulting from the energy loss the collisional and radiative energy loss contribute equally. Indeed, the quenching of hadron spectra observed at RHIC is significantly larger than that predicted solely from the radiative energy loss, which was accepted as a signature for the formation of a QGP phase in relativistic heavy-ion collisions at RHIC as the energy loss in hadronic matter is assumed to be smaller [29].

The dielectric functions (1) are derived from the Vlasov equation, i.e. for the case of a collisionless plasma. Of course, in most cases collisions play an important role in a plasma, e.g. for equilibration and transport properties, and cannot be neglected. Collisions can be considered by using the Boltzmann equation (Vlasov equation plus collision term). There are several approximations to this equation, of which the relaxation time approximation replacing the collision integral by a constant collision rate ν is the simplest one. In particular, the formalism proposed in [30] has been applied to low-temperature plasmas successfully

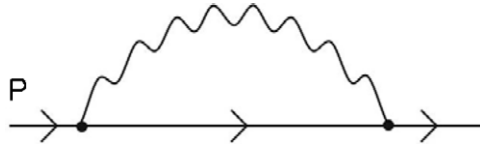


Figure 5. Fermion self-energy in lowest order perturbation theory.

(for complex plasmas, see e.g. [31]). In this case, an analytic expression for the dielectric functions containing the collision rate can be derived for ultra-relativistic plasmas [32]. For example, the longitudinal dielectric function now reads as

$$\epsilon_L(\omega, k) = 1 + \frac{3m_\gamma^2}{k^2} \left[1 - \frac{\omega + i\nu}{2k} \ln \frac{\omega + i\nu + k}{\omega + i\nu - k} \right] \left[1 - \frac{i\nu}{2k} \ln \frac{\omega + i\nu + k}{\omega + i\nu - k} \right]^{-1}. \quad (6)$$

The plasmon dispersion relation now cuts the light cone ($\omega = k$) at a finite value of k and the plasma frequency is given by $\omega_L(k = 0) < m_\gamma$. In the case of a QGP, m_γ is again replaced by m_g .

In heavy-ion collisions, the fireball expands dominantly in the longitudinal (beam) direction. Hence, the system is not isotropic and more components of the dielectric tensor have to be considered [33]. In anisotropically expanding plasmas, instabilities such as the Weibel instability show up. It has been shown that these instabilities can drive a QGP rapidly to equilibrium even in the weakly coupled case [34]. However, using the relaxation time approach as in (6) it was argued that these instabilities are suppressed to some extent by the presence of collisions in the plasma [35].

Ultra-relativistic plasmas offer the exciting possibility of observing a new kind of collective modes besides electromagnetic ones, namely fermionic modes. They cannot be derived from the dielectric function but from the electron or quark self-energy. In lowest order perturbation theory, the fermion self-energy is given by the diagram in figure 5. The pole of the effective fermion propagator following from a resummation of this self-energy describes the dispersion relation of collective fermion modes in the plasma. As in the case of electromagnetic (photon or gluon) waves, there are two branches: one with a positive ratio of helicity to chirality and one with a negative, which is absent in vacuum, called plasmino [36]. The plasmino branch shows a minimum at a finite value of the momentum p as sketched in figure 6. A minimum in the dispersion relation leads to van Hove singularities which could show up in sharp peaks in the dilepton production rate in the QGP, serving as a possible signature for the QGP formation [36]. It was argued that this minimum in the plasmino branch is a general feature of ultra-relativistic plasmas independent of the approximation, e.g. perturbation theory [37]. The observation of collective electron modes in an ultra-relativistic electron–positron plasma, produced by strong laser fields in the laboratory [38], would be an exciting discovery, also serving as a test for the QGP [39].

As a side remark, I would like to mention that the fermion self-energy can also be calculated perturbatively at high quark densities by introducing a finite chemical potential μ . In this way, for example, a quasi-particle mass for quarks in quark matter can be derived which describes the equation of state of a quasi-particle Fermi gas. The mass–radius relation for quark stars has been computed using this equation of state [40]. However, so far no indications for quark stars have been found (see e.g. [41]).

There are a number of other interesting quantities which can be calculated by perturbation theory going beyond the classical high-temperature approximation. For this purpose, one has to adopt the hard thermal loop resummation technique in order to obtain consistent,

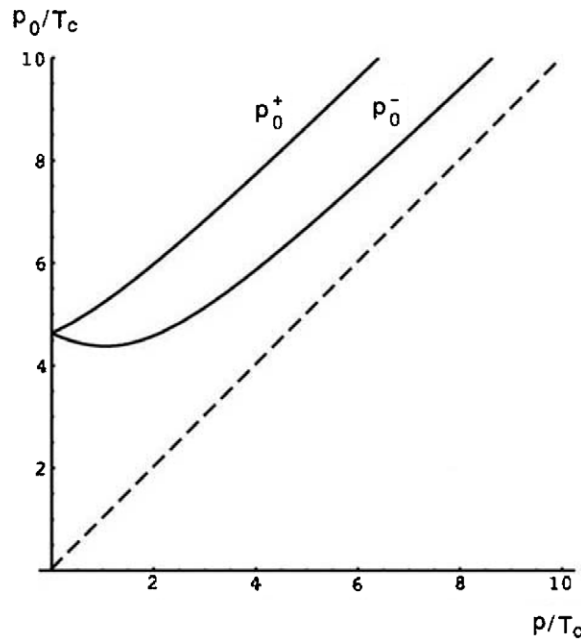


Figure 6. Dispersion relations of electrons in a QED plasma or quarks in a QGP.

i.e. complete to leading order, infrared finite, and gauge independent, results [26]. Important examples are damping rates, transport rates, mean free paths, collision times and transport coefficients (viscosity) of electrons, quarks, photons and gluons [15, 39]. Of course, these results hold only in the limit of extremely high temperatures—even at the Planck scale, the QCD coupling constant g is of the order of $1/2$. However, they might provide qualitative insight into important aspects of the physics of the QGP such as the role of collective effects. In the case of an electron–positron plasma, however, these calculations are reliable, serving as predictions for such a QED plasma which might be produced in strong laser fields soon [39].

3. The strongly coupled quark–gluon plasma

The essential parameter distinguishing between weakly coupled and strongly coupled non-relativistic plasmas is the Coulomb coupling parameter defined by the ratio of the interaction energy (Coulomb energy) between the plasma particles to their thermal energy [42]:

$$\Gamma = \frac{Q^2}{dk_B T}, \tag{7}$$

where Q is the charge of the particles, d is the interparticle distance and T is the plasma temperature. For strongly coupled plasmas, this parameter is of the order of 1 or larger. For example, for the ion component in a white dwarf Γ can be between about 5 and 500. In complex or dusty plasmas, which contain micron size particles, e.g. dust grains, the microparticle component is highly charged (several thousand elementary charges) due to electron collection and interacts via a Yukawa potential leading to a Γ between 1 and 10^5 depending on the

particle size and the plasma parameters [43]. A simple, extensively studied model for a strongly coupled plasma is the one-component plasma, in which particles of the same charge in a neutralizing background interact via a Coulomb potential [42]. For $\Gamma > 1$, the plasma shows a liquid-like behavior (see below), and for $\Gamma > 172$ a crystalline structure. Such a plasma crystal has been observed in complex plasmas in the laboratory [44]. An improvement of the one-component model is the Yukawa system taking into account the Debye screening of the particle charge.

In the case of a QGP, the interaction parameter was estimated to be [45]

$$\Gamma = 2 \frac{C \alpha_s}{dk_B T}, \quad (8)$$

where $C = 4/3$ is the Casimir invariant in the case of quarks and $C = 3$ in the case of gluons. The pre-factor of 2 has been added to consider the fact that in relativistic plasmas, the magnetic interaction is as important as the electric. Employing realistic values for RHIC energies, e.g. $T = 200$ MeV, $\alpha_s = 0.3\text{--}0.5$ and $d = 0.5$ fm, we find $\Gamma = 1.5\text{--}6$. Here, it should be noted that only the Coulomb potential corresponding to a one-gluon exchange was assumed. Higher order and non-perturbative effects can increase the value of Γ significantly. Anyway, this estimate indicates that the QGP in ultra-relativistic heavy-ion collisions is a strongly coupled plasma probably in the liquid phase. This means that there could be a phase transition to a QGP gas at higher temperatures where Γ will be smaller [46]. However, only in the simultaneous presence of attractive and repulsive interactions, such as a Lennard–Jones potential, this phase transition is of first order with a critical end point. Otherwise, the system is always in the supercritical phase allowing no determination of a sharp borderline between the liquid and the gaseous behavior.

The theoretical description of the strongly coupled QGP is difficult because perturbative QCD is not applicable and lattice QCD is restricted to static quantities and its accuracy at finite temperature—not to speak of finite chemical potential—not yet satisfactory in many cases. Therefore electromagnetic strongly coupled plasmas, which can be investigated more easily, and the methods, e.g. molecular dynamics [47, 48], for describing them are considered to improve our—at least qualitative—understanding of the QGP by analogy. For example, ultra-cold quantum gases exhibit a similar behavior in the flow pattern observed at RHIC [49]. The elliptic flow investigated in these heavy-ion collisions can be described very well by ideal hydrodynamics, indicating the presence of an almost ideal QGP liquid [50]. Also, high-density plasmas produced by shooting heavy-ion beams onto solid-state targets are another example of strongly coupled plasmas in the laboratory [51]. A model system for the QGP which can be produced for a wide range of values of Γ and investigated on the microscopic and dynamical level in real time by direct optical observation is the complex plasma, already discussed above. It has the further advantage that the strong coupling is due to the large coupling (high charge of the microparticles) as in the case of the QGP and not because of high density or low temperature.

Important tools for investigating fluids, such as complex plasmas in the liquid phase, on the microscopic level are the pair correlation function and its Fourier transform, the static structure function [52]. The qualitative behavior of the latter for the gas and liquid phases is shown in figure 7. Using the hard thermal loop approximation, valid in the weak coupling limit, a gas-like behavior of an interacting quark system is found [53]:

$$S(p) = \frac{2n_F T^3}{n} \frac{p^2}{p^2 + m_D^2}, \quad (9)$$

where n is the particle density, n_F is the number of quark flavors and $m_D = 1/\lambda_D$ is the inverse Debye screening length. It would be interesting to compute the static structure function

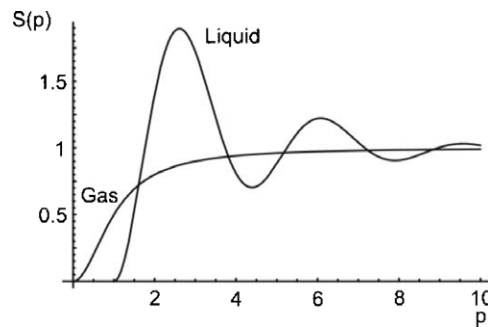


Figure 7. Qualitative behavior of the static structure function for the liquid and gaseous phases.

non-perturbatively for realistic situations, in particular by using lattice QCD, to see whether an oscillatory behavior indicating a QGP liquid will be found.

Another example of a lesson that can be learned from strongly coupled electromagnetic plasmas is the cross section enhancement of the particle scattering due to nonlinear effects. It has been argued that RHIC data imply a cross section enhancement in the elastic parton scattering by an order of magnitude compared to perturbative results [54]. Indeed in strongly coupled plasmas, such as complex plasmas [55], such a cross section enhancement takes place because the interaction range in these plasmas is larger than the Debye screening length which therefore cannot be used as an infrared cutoff in the calculation of the absolute cross section. An estimate of this effect in the QGP case gives a parton cross section enhancement by a factor of 2–9 implying a small mean free path, a small viscosity and a fast thermalization in accordance with RHIC experiments [45]. Furthermore, such a cross section enhancement would lead to a further increase of the collisional energy loss, given approximately by the energy transfer per collision divided by the mean free path. The radiative energy loss, on the other hand, could be suppressed due to an enhancement of the Landau–Pomeranchuk–Migdal effect, which describes a suppression of the photon or gluon radiation if the time between two scatterings is too small to allow the emission of the photon or gluon [46].

As a last example for an interesting comparison between the QGP and strongly coupled electromagnetic plasmas, let me mention the prediction of a lower limit for the ratio of viscosity to entropy density from string theory (AdS/CFT) [56]:

$$\frac{\eta}{s} \leq \frac{\hbar}{4\pi k_B}, \quad (10)$$

which is widely discussed in the QGP community because the QGP seems to come close to this limit [57]. In general, strongly interacting systems show a small viscosity, e.g. the one-component plasma has a minimum in the viscosity at $\Gamma = 21$ [58]. The minimum of the ratio of viscosity to entropy density in the one-component plasma at $\Gamma = 12$ is about five times above the string theory limit [59]—water under normal conditions exceeds this limit by about a factor of 400—similar to predictions for the QGP.

4. Conclusions

Transport theoretical methods (Vlasov and Boltzmann equation) widely used for non-relativistic electromagnetic plasmas and perturbative field theory at finite temperature (and chemical potential) can be used for describing the weakly coupled phase of the QGP,

e.g. collective and transport properties, and high-energy phenomena (jets, hard photons) in the QGP. Also, properties of a relativistic electron–positron plasma produced in ultra-strong laser fields or supernovae can be treated in this way. The comparison between these two systems might be helpful to learn about the role of the strong coupling effects in the QGP. Field theoretic predictions of a new phenomenon not known from non-relativistic plasmas, namely the existence of collective fermion modes (plasminos), might open exciting investigations of relativistic electron–positron plasmas in the laboratory.

Properties of a strongly coupled QGP such as a liquid phase, cross section enhancement and small viscosity can be studied in analogy to strongly coupled electromagnetic plasmas. In particular complex plasmas, which can be easily produced with a highly tunable interaction strength and directly investigated on the microscopic and dynamical level, showing a large variety of interesting features such as solid and liquid phases might be useful in this respect.

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