Comparing different hard-thermal-loop approaches to quark number susceptibilities

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Abstract. We compare our previously proposed hard-thermal-loop (HTL) resummed calculation of quark number susceptibilities using a self-consistent two-loop approximation to the quark density with a recent calculation of the same quantity at the one-loop level in a variant of HTL-screened perturbation theory. Besides pointing out conceptual problems with the latter approach, we show that it severely over-includes the leading-order interaction effects while including none of the plasmon term which after all is the reason to construct improved resummation schemes.

1 Introduction

In view of the ongoing search for quark-gluon plasma signals in the early stages of ultrarelativistic heavy-ion collisions, quark number susceptibilities (QNS) have recently received enhanced attention because of their direct connection with fluctuations of conserved charges which could in principle discriminate against a purely hadronic phase. Still, they give rise to a puzzling discrepancy in Ref. [7], implying a numerical value of \( \chi/\chi_0 \propto N_f \) for all temperatures \( T \lesssim 40T_c \), and only for \( T \gtrsim 700T_c \) does \( \chi/\chi_0 \) show the expected growth with temperature, starting from values extremely close to the ideal-gas result.

In Ref. [7] we have shown that these problems can be avoided by a reorganization of perturbation theory which is based on a self-consistent (\( \Phi \)-derivable) two-loop approximation to the thermodynamic potential \[ ] . The latter leads to a nonperturbative expression for entropy and quark density which can be used to resum the so-called hard thermal loops (HTL) \[ ] and particular next-to-leading order corrections thereof. The results for QNS thus obtained are monotonic functions of \( T/T_c \) which account at least for a sizeable part of the deviation from the ideal-gas behaviour observed in lattice calculations for \( T/T_c \gtrsim 2T_c \).

Recently, a different approach to resum the effects of HTL in QNS has been put forward in Ref. [11] which starts from quark number charge correlators. Employing HTL propagators and vertices at one-loop order, one finds substantially larger deviations from the ideal-gas limit, seemingly in a good agreement with the lattice results of Refs. [11].

In view of the large efforts invested at present by lattice gauge theorists to explore effects of small chemical potentials at high temperature in QCD, we think it worthwhile to explain the fundamental differences between our approach and that of Ref. [11] and why, in our opinion, the results of the latter are actually misleading. In particular we show that the one-loop results of Ref. [11] severely over-include the leading-order interaction effects, while they contain none of the plasmon effects \( \propto \alpha_s^{3/2} \) (which are the source of the problems with conventionally resummed perturbation theory). Moreover, we point out a certain...

\[ \frac{\chi}{\chi_0} = 1 - 2\frac{\alpha_s}{\pi} + 8\sqrt{1 + \frac{N_f}{6} \left( \frac{\alpha_s}{\pi} \right)^{3/2} + \left( \alpha_s^2 \log(\alpha_s) \right)^2} \]

\( N_f \) quark flavours. For \( N_f = 2 \) the plasmon term overcompensates the term \( \propto \alpha_s \) for all temperatures \( T \lesssim 40T_c \), and only for \( T \gtrsim 700T_c \) does \( \chi/\chi_0 \) show the expected growth with temperature, starting from values extremely close to the ideal-gas result.
technical difficulty that has been overlooked by the authors of Ref. [11], but has the effect to render their result ill-defined in a distributional sense.

More importantly even, we comment on a conceptual problem with the approach followed in Ref. [11], which arises because the HTL action is no longer used as the effective theory for soft modes, but is used throughout all of phase space. Implicitly the definition of the quark number charge operator is modified such as to no longer conform with the operator employed in lattice calculations.

Before discussing the approach of Ref. [11] in detail in Sect. 3, we briefly review the QNS as obtained from HTL-resummed thermodynamic potentials. Sect. 4 summarizes our conclusions.

2 QNS from resummed thermodynamic potentials

2.1 Generalities

The QNS of a given quark flavour is by definition the response of the quark number density $N$ to an infinitesimal variation of the associated chemical potential $\mu$,

$$\chi = \frac{\partial N}{\partial \mu} = \frac{\partial^2 P}{\partial \mu^2} = \beta \int d^3 x \langle \rho(0, x)\rho(0, 0) \rangle$$

(2)

where $P = (3\beta)^{-1} \log Z$ is the thermodynamic pressure, $\beta = T^{-1}$ and $\rho = \bar{\psi} \gamma^0 \psi$.

When thermodynamic consistency is automatic, for example in strict perturbation theory to a given order in $\alpha_s$, it does not matter which of the equivalent expressions on the right-hand side of (2) is employed. However, when further resumptions are performed that amount to a partial inclusion of higher-order effects, it does in fact matter. To set the stage we begin by briefly reviewing the approaches which focus on the thermodynamic potential before turning to Ref. [11] which starts from the quark number charge correlator.

Expressed as a functional of full propagators ($D$ for gauge bosons and $S$ for fermions, and assuming a ghost-free gauge choice) the thermodynamic potential $\Omega = -PV = -T \log Z$ has the form [12]

$$\beta \Omega[D, S] = \frac{1}{2} \text{Tr} \log D^{-1} - \frac{1}{2} \text{Tr} \Pi D - \text{Tr} \log S^{-1} + \text{Tr} \Sigma S + \Phi[D, S]$$

(3)

where $\Phi$ is the sum of 2-particle-irreducible “skeleton” diagrams whose lowest-order (2-loop) contributions are

$$\Phi[D, S] = -1/12 + 1/8 + 1/2 + 1/8 + \cdots$$

As a functional of $D$ and $S$, $\Omega$ is subject to the stationarity condition,

$$\delta \Omega[D, S] / \delta D = 0 = \delta \Omega[D, S] / \delta S,$$

(4)

which is equivalent to

$$\delta \Phi[D, S] / \delta D = \frac{1}{2} \Pi, \quad \delta \Phi[D, S] / \delta S = \Sigma,$$

(5)

for the self-energies $\Pi$ and $\Sigma$. Expressing $\Pi = D^{-1} - D_0^{-1}$ and $\Sigma = S^{-1} - S_0^{-1}$ in terms of bare propagators $D_0$ and $S_0$, the representation [11] of course reproduces the ordinary loop expansion.

For example, the leading-order interaction terms $\propto \alpha_s$ are given by the 2-loop diagrams in $\Phi$, whereas single powers of the self-energy insertions in a propagator cancel out in the first four terms of the right-hand side of eq. (3).

Ordinary perturbation theory, however, has infrared problems at finite temperature if the repeated self-energy insertions contained in the term $1/2 \text{Tr} \log D^{-1}$ are expanded out perturbatively. These can be remedied by a resummation of the leading order Debye mass

$$m_D^2 = (2N + N_f) \frac{g^2 T^2}{6} + \sum_i \frac{g^2 \mu_i^2}{2\pi^2}$$

(6)

in the (chromo-)electrostatic propagator, where $g^2 = 4\pi \alpha_s$ (though new infrared problems arise at order $\alpha_s^3$). Expanded in powers of $g$, the resummation of the self-energy term in the pressure $P_3 = N_f T m_D^2 / (12\pi)$.

It is this term which is responsible for the dramatic deterioration of the apparent convergence of a perturbative expansion of $P$ in $g$ at finite temperature, and, as remarked in the introduction, to a somewhat lesser degree for QNS which can be derived from the pressure.

2.2 Screened (HTL) perturbation theory

The loss of apparent convergence upon inclusion of the plasmon term in the pressure is in fact generic and also occurs in a simple scalar $\varphi^4$ theory [13]. This problem arises as soon as finite-temperature contributions are expanded out in powers of the coupling, which is necessary for the standard ultraviolet renormalization programme to become applicable. In order to avoid this, it has been proposed [14] to reorganize perturbation theory by adding screening masses to the classical Lagrangian and to subtract them as counter-terms, but in contrast to the usual resummation programme at finite temperature [11,15], this is done for both hard and soft momentum regimes. This in fact alters the ultraviolet structure of the theory, but when combined with a simple minimal subtraction of the additional divergences this resummation appears to significantly improve the apparent convergence of thermal perturbation theory.

In Refs. [13,14] this approach has been extended to QCD at one-loop level. It amounts to keeping only the logarithmic terms in (3) and replacing $D$ and $S$ by the HTL propagators,

$$\beta \Omega^{1-loop-HTL} = \frac{1}{2} \text{Tr} \log \tilde{D}^{-1} - \text{Tr} \log \tilde{S}^{-1}$$

(8)
where hatted quantities refer to HTL. If the thermal mass parameter in the HTL are exactly the lowest-order ones, this includes the correct plasmon term \( \propto \alpha_s \) without causing the pressure to exceed the ideal-gas value. However, the leading-order interaction pressure \( \propto \alpha_s \) is over-included by a factor of 2\(^2\).

In order to have both, the leading-order term \( \propto \alpha_s \) and the plasmon term \( \alpha_s^{3/2} \) included correctly, it is necessary to go to two-loop order. Starting from two-loop order, one can turn this so-called “HTL perturbation theory” into a variational perturbation theory, where the HTL action is no longer used as an effective theory for soft modes as in standard HTL resummation\(^2\), but just as a gauge invariant mass term which is then eliminated by a principle of minimal sensitivity.

Because of the HTL involves non-local self-energies and vertices, this optimization of perturbation theory is extremely difficult and has only recently been carried through for QCD to 2-loop order\(^3\). The results are a clear improvement over conventionally resummed perturbation theory, where the HTL action is used uniformly for soft and hard momenta, whereas the HTL action is accurate only for soft momenta, and for hard ones only in the vicinity of the light-cone. A related problem is that the artificial UV divergences that are introduced involve new subtraction scheme dependences. While these start to be suppressed by powers of \( \alpha_s \) only at the (rather forbidding) three-loop order\(^3\), these additional scheme dependences turn out to be numerically rather weak in the two-loop result for QCD.

2.3 HTL-resummation of the 2-loop \( \Phi \)-derivable entropy and density

While HTL-screened perturbation is in principle rather generally applicable, we have found, following up an observation made in Ref.\(^3\), that specifically for the first derivatives of the thermodynamic potential one can derive remarkably simple expressions from a self-consistent 2-loop approximation to the skeleton expansion\(^3\) of the QCD thermodynamic potential. Because of the stationarity property, these derivatives act only on the explicit statistical distribution functions, and not also on those contained in propagators and self-energies. Moreover, after differentiation, the contribution from \( \Phi^{2-\text{loop}} \) just cancels part of the second and fourth term on the right-hand-side of (6). The derivatives with respect to temperature and chemical potential give entropy and quark densities, respectively, reading

\[
S = -\text{tr} \int \frac{d^4k}{(2\pi)^4} \frac{\partial n}{\partial T} \left[ \text{Im} \log D^{-1} - \text{Im} \Sigma \text{Re} D \right] -2 \text{tr} \int \frac{d^4k}{(2\pi)^4} \frac{\partial f}{\partial T} \left[ \text{Im} \log S^{-1} - \text{Im} \Sigma \text{Re} S \right],
\]

\[
\mathcal{N} = -2 \text{tr} \int \frac{d^4k}{(2\pi)^4} \frac{\partial f}{\partial \mu} \left[ \text{Im} \log S^{-1} - \text{Im} \Sigma \text{Re} S \right],
\]

where \( D \) and \( S \) are determined by one-loop gap equations\(^3\) obtained by restricting \( \Phi \) to two-loop order.

Although these equations have the form of one-loop expressions involving dressed propagators, they include all the two-loop contributions, but incorporated in the spectral properties of the quasi-particles described by the dressed propagators. This implies that both the leading-order interaction terms \( \propto \alpha_s \) and the plasmon effect \( \propto \alpha_s^{3/2} \) are completely taken into account as soon as \( D \) and \( S \) are evaluated to sufficient accuracy.

Because the expressions (9) and (10) are manifestly ultraviolet finite, they can be used to resum the effects of HTL without the necessity of subsequent expansions and truncations (see Ref.\(^3\)). At soft momenta, the HTL are valid expressions to the actual full propagators that one would have to use in a self-consistent scheme; at hard momenta it turns out that to leading order the above expressions only probe the vicinity of the light-cone where HTL self-energies remain accurate. The next-to-leading order effect, which is the plasmon effect, turns out to be to one part covered by HTL resummation in the soft regime, and to the remaining part by corrections to the so-called asymptotic thermal masses from HTL-resummed one-loop diagrams.

In the case of the quark number density functional, from which the QNS can be derived, it turns out that all of the plasmon effect is associated with next-to-leading order corrections to the asymptotic fermion mass, shown in Fig. (1).

If only the HTL approximation to the fermion propagator is employed, the quark number density still contains the complete leading-order interaction effects \( \propto \alpha_s \), and – since it need not be expanded out perturbatively – also subsets of higher-order effects, namely those associated with repeated HTL insertions.

In Ref.\(^3\) we have evaluated the QNS obtained by taking the derivative of (11) with respect to \( \mu \), both in the HTL approximation and in a next-to-leading approximation which incorporates the plasmon effect through the corrections to the asymptotic fermion mass from the diagram of Fig. (1).

\(^2\) As explained in the last paper of Ref.\(^3\), Ref.\(^3\) had an even stronger over-inclusion due to an inconsistent use of dimensional regularization.

\(^3\) For a different approach based on the pressure see also Ref.\(^3\).
3 QNS from HTL-resummed charge correlators

In Ref. [11] an HTL-resummed QNS has been constructed by starting from the charge correlator

\[ \chi = \beta \int d^3x \langle \rho(0,\mathbf{x})\rho(0,\mathbf{0}) \rangle = \beta \int d^3x \Pi^>_{\mu\nu}(0,\mathbf{x}) \]  

(11)

where \( \Pi^>_{\mu\nu} \) is the current-current correlator of a given flavour charge (suitable linear combinations of such quantities give the correlators of electric charge and baryon number).

In Fourier-space one has \( \chi = \lim_{k \to 0} \beta \int_{-\infty}^{\infty} \frac{d\omega}{2\pi} \Pi^>_{\mu\nu}(\omega, k) \)

(12)

where the well-known fluctuation-dissipation theorem allows one to write (in the notation of [11])

\[ \Pi^>_{\mu\nu}(\omega, k) = -\frac{2}{1 - e^{-\beta \omega}} \text{Im} \Pi^{R\mu\nu}(\omega, k) \]

(13)

with \( \Pi^{R\mu\nu} \) the retarded response function.

In conventional perturbation theory the first few diagrams contained in (12) are shown in Fig. 2 [21,22,23]. The first diagram gives the ideal gas value, and the leading-order interaction term \( \propto \alpha_s \) is given by the two-loop order diagrams. The plasmon effect \( \propto \alpha_s^{3/2} \) comes from those higher-order diagrams which correspond to repeated self-energy insertions into the gluon lines of the two-loop diagrams of Fig. 3.

Calculating \( \chi \) from (12) is in fact a bit more involved than starting from the thermodynamic potential as also noticed in Ref. [21]. Because charge conservation implies \( \omega \Pi^>_{\mu\nu}(\omega, 0) = 0 \) one has \( \lim_{k \to 0} \Pi^>_{\mu\nu}(\omega, k) \sim \delta(\omega) \).

If, for example following HTL-screened perturbation theory, one is interested in the one-loop contribution arising from dressing the propagators in Fig. 2, this will spoil this behaviour. In order to have charge conservation one needs to employ HTL vertices in addition to HTL propagators, which is precisely what the authors of Ref. [11] have proposed.

However, this raises an important conceptional problem: in effect this use of the HTL vertices replaces the ordinary charge operator \( \psi \gamma^0 \psi \) by the non-local object \( \psi(\gamma^0 + F^0) \psi \) derivable from the non-local HTL action. The correspondingly re-defined QNS is therefore no longer directly related to the quantity defined in (11) and measured in lattice simulations.

This is in fact a problem only because the HTL action is no longer used as an effective theory appropriate for soft momentum scales, but is used equally for soft and hard momenta. As an effective theory, obtained after integrating out the hard momenta and used for soft modes only, the appropriate charge operator is indeed the non-local quantity involving the HTL vertex \( F^0 \). It is this operator which enters in a perturbative matching to the full theory. But replacing the ordinary charge operator by the HTL-dressed one for all momenta clearly corresponds to abandoning the definition (11) for the QNS.

Leaving this issue aside for now, we continue by analysing the diagrammatic content of \( \chi \) in HTL-screened perturbation theory. Since HTL vertices are strictly one-loop quantities, to one-loop order it is as shown in Fig. 3 [21].

However, while all the topologies that are present in the two-loop diagrams of Fig. 3 also appear in the diagrams of Fig. 3, their combinatorial factors are different. This shows that in the one-loop HTL approximation to (12) the \( \alpha_s \)-contributions are all overcounted. This is the second diagram of Fig. 3 is contained with correct combinatorics in the first diagram of the right-hand side of Fig. 3 but appears another time through the HTL 4-vertex; the third diagram of Fig. 3 is seen to be over-included by a factor of 2. Moreover, because the HTL approximation for the undressed self-energies and vertex subdiagrams in Fig. 3 does not provide the complete leading-order terms for hard inflowing momenta, this leads to a further source of incompleteness of the terms \( \propto \alpha_s \).

The authors of Ref. [11] do not specify how to extend their approach to two-loop order. It is clear, however, that the correct counting is restored only when the modification of the quark number charge is undone by HTL counter-terms and all two-loop diagrams are added. The relevant diagrams for completing the order \( \alpha_s \) result are shown in Fig. 4, where the first and the third diagram...
correspond to HTL counter-terms to the charge operator. This means that the definition of the charge operator in Ref. [11] has to be modified order by order to approach to standard definition of QNS at least at infinite loop order.

This also shows that the plasmon term \( \propto \alpha_s^{3/2} \), which is the reason for seeking improvements of conventionally resummed perturbation theory, only appears in the two-loop order diagrams of HTL-screened perturbation theory, namely through the dressed vector boson lines with a blob. The vector boson lines within the HTL vertices of Fig. 3 are not dressed and thus do not capture anything of the plasmon effect.

We are now in a position to compare with the HTL-resummation approaches discussed in the previous section.

In one-loop HTL-screened perturbation theory along the lines of Refs. [15,16] the leading-order interaction term to the thermodynamic potential is over-included by a factor 2, but the plasmon effect is complete (as long as only the leading-order HTL mass parameter is used; including higher-order corrections in the latter would spoil this). The QNS have not been calculated in this approach, but the leading-order HTL mass parameter is used; including higher-order corrections in the latter would spoil this).

The 2nd and 4th diagram have opposite combinatorial factors but may also contribute to order \( \alpha_s \), because of the incompleteness of the HTL approximation for hard loop momenta. Giving up thermodynamic consistency, one could think of identifying the mass parameter in HTL-screened perturbation theory only after this differentiation. This would in fact give a correct leading-order interaction term, but would loose the plasmon effect (and not completely reproduce the terms of order \( \alpha_s^2 \) and higher contained in Ref. [12]).

In Figs. 5 and 6 we consider the expression

\[ R \equiv (\chi - \chi_0)/(\chi_{p.th}^{(2)} - \chi_0) \]

(14)

which measures the deviation of the interaction part of \( \chi \) from the perturbative result \( \chi_{p.th}^{(2)} \) to order \( \alpha_s \), and plot the respective results, again as a function of \( M/T \). While our result shows slightly slower deviation from the ideal-gas limit than the strictly perturbative result to order \( \alpha_s^2 \), the result of Ref. [12] has considerably stronger deviations because of the over-inclusion of the leading-order interaction term.

In Fig. 5 we consider the expression

\[ \frac{\chi}{\chi_0} \quad \text{as a function of} \quad \frac{M}{T} \]

for the 1-loop HTL result of Ref. [11] (upper half of the plot) and for the 2-loop \( \Phi \)-derivable approximation evaluated with HTL propagators (lower half of the plot — note the different scale!). The perturbative result to order \( g^2 \) corresponds to the value 1. Only the 2-loop HTL result approaches 1 in the limit \( M/T \to 0 \), where eventually perturbation theory should be reproduced; the 1-loop HTL result of Ref. [11] is seen to over-include the leading-order interaction effect by a factor which diverges logarithmically as \( M/T \propto g \to 0 \).
This clearly shows that the one-loop HTL resummation of the charge-charge correlator cannot be compared with either our results, which are based on the two-loop expression (11), or perturbation theory which it seeks to improve upon. In our opinion it is also completely premature to compare with the available lattice results on QNS, because the two-loop contributions of Fig. 4 will have to correct for the enormous over-inclusion of terms $\propto \alpha_s$.

But it appears to be questionable whether a two-loop HTL-screened perturbation theory calculation of (12) is at all practicable. There are in fact certain technical problems with the result reported in Ref. [11] already at one-loop order. In eq. (34) of [11] one can see that the result for $\Pi_{00}^>(\omega, 0)$ is proportional to the integral

$$\int dx \int d^3x' n_F(x) n_F(x') \rho_+(x, k) \rho_-(x, k)$$

$$\times \frac{\left(\omega - x - x'\right)^2 \delta\left(\omega - x - x'\right)}{\omega^2}$$

(15)

where $\rho_\pm(x, k)$ are the HTL spectral functions for the two fermionic quasi-particle branches of the HTL approximation. In Ref. [11] the latter are used to put $x = -x'$, so that the second line of (15) is reduced to $\delta(\omega)$ in conformity with the expectations from charge conservation. However, this term is clearly ill-defined and might with equal justification be put to zero identically as is suggested by the way we have written it. In order to have a well-defined expression, it seems to be necessary to keep the external spatial momentum different from zero and take the limit to zero only after having performed the integral over $\omega$ (as demanded by (12)). But that would make its evaluation in HTL perturbation theory a hopelessly difficult task, already at one-loop order.

### 4 Conclusions

In this paper, we have discussed various possibilities for HTL-resummation in the calculation of QNS and have in particular analysed the recent proposal of Ref. [11]. We have shown that the resummed one-loop calculation presented there severely over-includes the leading-order interaction terms, while not including anything of the plasma effect, both of which would be corrected only at two-loop order. Thus only the latter should be viewed as an improvement over ordinary perturbation theory and used as such in a comparison with the available lattice results.

However, because of the technical problems mentioned at the end of the previous section, it would seem to be more sensible to calculate the QNS through a 2-loop HTL-screened perturbation theory evaluation of the pressure along the lines of Ref. [7].

On the other hand, the HTL-resummed calculation of QNS of Ref. [8] is based on a 2-loop $\Phi$-derivative approximation and does include correctly both the leading-order interaction effect, $\sim \alpha_s$, and the next-to-leading-order one, $\sim \alpha_s^{3/2}$ (together with an infinite series of higher order effects due to HTL). The results in Ref. [8] show the same trend as the lattice results, but a significant difference still remains, which calls for further studies, both on the analytic side, by further improving the resummation schemes, and on the lattice side, by increasing the reliability of the numerical results.

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