Consistent HTL resummation of the thermodynamical potential

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Abstract

The thermodynamical potential of relativistic plasmas with gauge interaction can be consistently resummed in terms of HTL propagators, which is, without being restricted to it, exemplified for the case of hot QED. The nonperturbative resummation obtained in a $\Phi$-derivable approach is gauge independent, free of thermal divergences and, in the weak-coupling limit, compatible with the leading order perturbative result.
I. INTRODUCTION

One of the main issues of the ongoing heavy-ion program is the investigation of deconfined hadronic matter. Despite asymptotic freedom of the strong interaction, this quark-gluon plasma is characterized by a large coupling in the regimes of physical interest, so nonperturbative approaches are required to describe this many-particle system reliably.

Recently, within the framework of an equilibrium description of the QCD plasma, the calculation of thermodynamical quantities by resumming hard-thermal-loop (HTL) propagators was proposed. In [1] and [2], nonperturbative expressions were given for the thermodynamical potential which, however, are detracted from inconsistencies: The expressions do not reproduce, in the weak-coupling limit, the perturbative results at leading order \( O(g^2) \) and they are, with uncompensated medium-dependent divergences, not satisfactory from a formal point of view. These problems were claimed to be solved only in a fully resummed calculation. In Ref. [3], instead, a consistent approximation of the entropy was derived from HTL propagators. This approach in principle resolves the problem of a leading order thermodynamical resummation since, up to an integration constant, the thermodynamical potential can be reconstructed from the entropy [4].

Nonetheless, a direct calculation of the thermodynamical potential, the quantity containing the full thermodynamical information, is instructive, in particular since an approach to derive macroscopic properties, which are sensitive to hard momenta, from HTL propagators (as a soft-momentum approximation) is, a priori, far from obvious. On the other hand, being explicitly gauge independent and respecting for arbitrary momenta the fundamental sum rules resulting from the commutator relations, the HTL propagators are a favorable basis for calculating physical quantities from approximate dressed Green’s functions, yielding (mostly) analytical nonperturbative results. These points are addressed in the present note where, starting from the Luttinger-Ward representation of the thermodynamical potential [5], the hot QED plasma is studied. This case is particularly simple to analyze but at the same time also representative for systems with an equal HTL structure as, e.g., the
II. THE APPROXIMATION SCHEME

As an exact relation, the thermodynamical potential can be expressed in terms of fully dressed Green’s functions by the (generalized) Luttinger-Ward representation \[5–7\]

\[
\Omega = \frac{1}{2} \tilde{\Omega}[D] + \tilde{\Omega}[S] + \Phi[D, S].
\]

(1)

\(\tilde{\Omega}\) is a functional of the photon propagator \(D\) or the electron propagator \(S\), which are related by Dyson’s equation to the respective self-energies. The boson part, e.g., is defined by

\[
\frac{1}{2} \tilde{\Omega}[D] = \frac{1}{2} \text{Tr} \left[ \ln(-D^{-1}) + D\Pi \right], \quad D^{-1} = D_0^{-1} - \Pi,
\]

where the trace is taken over the four-momentum and the Lorentz structure, while in the analog fermion part \(-\tilde{\Omega}[S]\), with \(S^{-1} = S_0^{-1} + \Sigma\), the spinor indices are traced. The functional \(\Phi[D, S]\) given by all two-particle irreducible bubble graphs with exact propagators (‘dressed skeletons’) is related to the self-energies by

\[
\Pi = -2 \frac{\delta \Phi}{\delta D}, \quad \Sigma = \frac{\delta \Phi}{\delta S}.
\]

(2)

Consequently, the fundamental stationarity of the thermodynamical potential upon variation of the self-energies \[8\] is fulfilled by the representation (1). It is emphasized that, on account of its stationarity, \(\Omega\) is formally gauge independent, although the propagators are not.

Using the projectors \(P^L_{\mu\nu} = -\tilde{K}_\mu \tilde{K}_\nu / K^2\) and \(P^T_{\mu\nu} = g_{\mu\nu} - K_\mu K_\nu / K^2 - P^L_{\mu\nu}\), where \(\tilde{K} = [K(Ku) - uK^2]/[(Ku)^2 - K^2]^{1/2}\) and \(u\) is the medium four-velocity, the inverse photon propagator is decomposed into the transverse (T) and the longitudinal (L) part as well as the covariant gauge-fixing term,

\[
D^{-1}_{\mu\nu}(K) = \sum_{i=T,L} P^i_{\mu\nu} \Delta^{-1}_i + \frac{1}{\xi} K_\mu K_\nu, \quad \Delta^{-1}_i = \Delta^{-1}_0 - \Pi_i,
\]

with \(\Delta^{-1}_0 = K^2 = k^2_0 - k^2\). Introducing the fermionic ‘projectors’ \(P_{\pm}(K) = \frac{1}{2} (\tilde{K} \pm \tilde{K})\) on the particle and the hole excitations (the index denotes the ratio of chirality to helicity), the electron propagator can be written in a similar way as
\[ S = \sum_{i=\pm} P_i \Delta_i, \quad \Delta_i^{-1} = \Delta_0^{-1} - \Sigma_i. \]

In terms of the scalar propagators \( \Delta_i \), with the degeneracy factors \( d_T = d - 1, d_L = 1 \) and \( d_\pm = (d + 1)/2 \), the \( \tilde{\Omega} \) parts of eq. (1) read

\[
\frac{1}{2} \tilde{\Omega}[D] = \frac{1}{2} \sum_{i=T,L} \left\{ \sum_i d_i \left[ \ln(-\Delta_i^{-1}) + \Delta_i \Pi_i \right] - \ln(-\Delta_0^{-1}) \right\},
\]

\[
\tilde{\Omega}[S] = \sum_{i=\pm} \left\{ \sum_i d_i \left[ \ln(-\Delta_i^{-1}) + \Delta_i \Delta_0^{-1} \right] - d_\pm \ln(-\Delta_0^{-1}) \right\},
\]

where the subtractive contribution of the ghost fields, which otherwise decouple, is included in \( \tilde{\Omega}[D] \), and the integral-sums, continued to \( d = 3 - 2\varepsilon \) spatial dimensions,

\[
\sum_{k^d} = \int_{k_0} T \sum_{k_0}, \quad \int_{k^d} = \mu^{2\varepsilon} \int \frac{d^d k}{(2\pi)^d},
\]

run over either bosonic or fermionic Matsubara frequencies \( k_0 \).

Commencing from this exact approach, selfconsistent (‘symmetry conserving’) approximations can be derived [9]; an approximation of the functional \( \Phi \) in the scheme (1-3) yields approximate self-energies and an expression for the thermodynamical potential which is still stationary. In particular, the loop expansion of \( \Phi \) can be truncated at a certain order, which is, in terms of the perturbative expansion in free Green’s functions, equivalent to a partial resummation avoiding the problem of double counting of diagrams. The leading-loop order of the \( \Phi \)-derivable approximation is given diagrammatically by [2]

\[ \Phi^{ll} = -\frac{1}{2} \bigcirc \bigcirc, \quad \Pi^{ll} = \sim \sim, \quad \Sigma^{ll} = - \bigcirc \bigcirc, \]

which implies the relation

\[ \Phi^{ll} = -\frac{1}{2} \Tr D^{ll} \Pi^{ll} = \frac{1}{2} \Tr S^{ll} \Sigma^{ll} \]

among the selfconsistent solutions of the coupled Dyson equations and \( \Phi^{ll} \).

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1 It is noted that the (vacuum) counter terms required to render the self-energies finite can be implemented in the \( \Phi \)-derivable scheme by appropriate counter loops in \( \Phi \), leaving \( \Omega \) itself unchanged.
In the following, the selfconsistent one-loop self-energies are approximated by the HTL self-energies $\Pi^*$ and $\Sigma^*$. Although the HTL self-energies remain a reasonable approximation even for hard momenta \([10,11]\), it is not evident a priori that a consistent approximation of $\Omega$ can be formulated in terms of such approximated quantities which are derived for soft momenta much smaller than the temperature $T$, whereas thermodynamics is sensitive to the momentum scale $T$. In fact, the HTL approximation of the self-energies undermines the selfconsistency of the $\Phi$-derivable approach since the relation analogous to (4) is not fulfilled. Nevertheless, as shown in the following, it is possible to consistently resum the HTL contributions to the thermodynamical potential, yielding a nonperturbative ‘continuation’ of the $O(e^2)$ perturbative result. In Sec. III, the HTL contributions to the $\tilde{\Omega}$ parts of the thermodynamical potential are calculated directly from the HTL self-energies since, e.g., replacing $\Pi^{\mu\nu}$ by $\Pi^*$ is correct up to terms beyond the order $O(e^2)$ under consideration. The corresponding $\Phi$ contribution is then given in Sec. IV to complete the resummed approximation of the thermodynamical potential which is, expressed by HTL self-energies, explicitly gauge invariant.

III. THE $\tilde{\Omega}$ CONTRIBUTIONS

The HTL self-energies of the photon and the electron are given by \([12]\)

$$
\begin{align*}
\Pi_T^* &= M_b^2 + \Pi, & \Pi_L^* &= -2\Pi, & \Pi(k_0, k) &= M_b^2 \frac{K^2}{k^2} \left[1 + \frac{k_0}{2k} \ln \frac{k_0 - k}{k_0 + k}\right], \\
\Sigma_{\pm}^* &= \frac{1}{2} M_f^2 \pm \tilde{\Sigma}, & \tilde{\Sigma}(k_0, k) &= \frac{M_f^2}{2} \left[\frac{k_0}{k} + \frac{K^2}{2k^2} \ln \frac{k_0 - k}{k_0 + k}\right].
\end{align*}
$$

The quantities

$$
M_b^2 = \frac{e^2 T^2}{6}, \quad M_f^2 = \frac{e^2 T^2}{4}
$$

can be considered as asymptotic masses (squared) of the transverse photon and the electron particle excitation, respectively, since their dispersion relations approach mass shells for momenta $k \gg eT$. The longitudinal photon (plasmon) mode and the hole (plasmino)
excitation, on the other hand, possess a vanishing spectral strength when approaching the light cone exponentially fast for $k \sim eT$.

Besides the $\Phi$ part, eq. (3) evaluated with HTL propagators yields the contributions to the thermodynamical potential which are, up to gauge group factors, the same as for non-Abelian gauge theories with HTL self-energies with a structure like (5). In Refs. [1,2], however, where the QCD plasma was studied in the framework of the leading order HTL perturbation theory, only the $\ln(-\Delta^{-1})$ terms of (3) were considered to contribute to the thermodynamical potential. In terms of the bare-propagator expansion, this incomplete analysis amounts to a miscounting of graphs, so in the weak-coupling expansion the leading order perturbative contribution to the thermodynamical potential is not reproduced correctly in either the case of hot QCD with vanishing chemical potential [1] or for the degenerate quark-gluon plasma [2] at $T = 0$. For the QED plasma under consideration, the corresponding boson ($b$) and fermion ($f$) contributions, marked here by the index $A$, follow from the expressions given in [1] by replacing the asymptotic gluon and quark masses by the expressions (6), schematically

$$\tilde{\Omega}_{A,b} = \frac{1}{2} \sum \left\{ d_T \ln(-\Delta_T^{-1}) + \ln(-\Delta_L^{-1}) - \ln(-\Delta_0^{-1}) \right\} = -\frac{M_b^4}{32\pi^2\varepsilon} + \text{finite terms},$$

$$\tilde{\Omega}_{A,f} = \sum d_{\pm} \left\{ \ln(-\Delta_+^{-1}) + \ln(-\Delta_-^{-1}) - \ln(-\Delta_0^{-1}) \right\} = \text{finite terms}.$$  \hspace{1cm} (7)

The imperfect cancelation of the thermal divergences in $\tilde{\Omega}_{A,b}$ is a second, formal indication of missing contributions in the approaches [1,2].

The complete analysis of the HTL-resummed thermodynamical potential has to keep track of the remaining terms, indexed $B$ in the following, of eq. (3). Using complex contour integration, the Matsubara sums can be calculated to yield a quasiparticle part stemming from the pole $\omega_i(k)$ of the propagators, and a Landau-damping contribution arising from the discontinuity of the HTL self-energies below the light cone,

$$\sum_{k_0} \Delta_i^* f_i = -(1 + 2n(\omega/T)) \left. \frac{f_i}{\partial_{\omega} \Delta_i^{-1}} \right|_{\omega_i} + \int_0^k \frac{d\omega}{2\pi} (1 + 2n(\omega/T))\Psi_i, \quad \Psi_i = \text{Disc}(\Delta_i^* f_i),$$

where $n = \pm n_b,f$ is either the Bose or the negative of the Fermi distribution function, and $f_i$ can be $\Pi_{T,L}^*$ or $\Delta_0^{-1}$. Subtracting and adding the appropriately infrared-regularized
asymptotic integrands analog to the technique applied in [1], the integrals over the spatial momenta are split into a finite contribution and the dimensionally regularized subtraction term. Taking the limit $\varepsilon \to 0$, the individual contributions are

$$\tilde{\Omega}_{B,b}^{\ast} = \frac{1}{2} \sum_{i=T,L} \sum_{d_i} \int_{k^3} \left[ \frac{-(1 + 2 n_b) \Pi_i^*}{2 \omega - \partial_\omega \Pi_i^*} \right]_{\omega_i} + \int_0^k \frac{d\omega}{2\pi} (1 + 2 n_b) \Psi_i - B_i^{\text{sub}} \right]$$

$$+ \frac{M_b^4}{32\pi^2} \left[ \frac{2}{\varepsilon} + 2 \ln \frac{4\pi}{e\gamma} + \frac{1}{3} \left( \frac{14}{3} - 2\pi^2 + \frac{16}{3} \ln 2 - \frac{8}{3} \ln 2 \right) \right] \left( \frac{M_b^2}{\mu^2} \right)^{-\varepsilon},$$

$$\tilde{\Omega}_{B,f}^{\ast} = \sum_{i=\pm} d_i \sum_{\Delta_i} \Delta_i \Delta_0^{-1} = \sum_{i=\pm} \int_{k^3} \left[ \frac{-(1 - 2 n_f) K^2}{2 \omega - \partial_\omega \Sigma_i^*} \right]_{\omega_i} + \int_0^k \frac{d\omega}{2\pi} (1 - 2 n_f) \Psi_i - B_i^{\text{sub}} \right]$$

$$+ \frac{M_f^4}{8\pi^2} \ln 2 \left( -1 + 2 \ln 2 \right),$$

with the angles $\Psi_{T,L} = \text{Disc}(\Delta_{T,L}^* \Pi_{T,L}^*), \Psi_{\pm} = \text{Disc}(\Delta_i^* \Delta_0^{-1})$, and the subtraction terms

$$\sum_{i=T,L} d_i B_i^{\text{sub}} = -d_T \left( \frac{M_b^2}{2k} + \frac{M_b^4}{2k(k^2 + M_b^2)} \left( 2 - \ln \frac{4(k^2 + M_b^2)}{M_b^2} \right) \right)$$

$$- \int_0^k \frac{d\omega}{\pi} \text{Im} \tilde{\Pi} \left[ \frac{-d_T K^2}{(K^2 - M_b^2)^2} \left( 1 + 2 \frac{\text{Re} \tilde{\Pi}}{K^2 - M_b^2} \right) + \frac{2}{K^2} \left( 1 - 4 \frac{\text{Re} \tilde{\Pi}}{K^2} \frac{k^2}{k^2 + M_b^2} \right) \right],$$

$$\sum_{i=\pm} d_i B_i^{\text{sub}} = -\frac{M_f^2}{2k} - \frac{M_f^4}{8k(k^2 + M_f^2)} \left( 1 - 2 \ln \frac{4(k^2 + M_f^2)}{M_f^2} \right)$$

$$- \int_0^k \frac{d\omega}{\pi} \text{Im} \tilde{\Sigma} \left[ \frac{1}{K^2} - \frac{K^2}{K^2 - M_f^2} \right].$$

It turns out in eq. (8) that the thermal divergences of the fermion integral-sum cancel, as for $\tilde{\Omega}_{A,f}^{\ast}$ in (7). The boson part, on the other hand, contains a temperature dependent term $\sim M_b^4/\varepsilon$ which is twice as large in magnitude as its counterpart in (7). These thermal divergences of the $A$ and $B$ contribution have to cancel the corresponding terms in the HTL approximation of $\Phi$, which is calculated in the following section, to yield a well-defined resummation of the thermodynamical potential.
IV. THE $\Phi$ CONTRIBUTION

In this section, the remaining $\Phi$ contribution is calculated by evaluating the two-loop functional, given the self-energies and the Green’s functions. Within the leading loop approximation, $\Phi$ is related to the self-energies by equation (4) and can thus be expressed in a general form by $\Phi^\mu = -\frac{1}{2} t \text{Tr} D^\mu \Pi^\mu + \frac{1}{2} (1 - t) \text{Tr} S^\mu \Sigma^\mu$, independently of $t$. In the framework of the HTL approximation, however, the consistency relation (4) is violated since the traces are dominated by hard momenta. The resulting ambiguity is even of order $\mathcal{O}(e^2)$, namely

$$\begin{align*}
\left. \Phi^\mu \right|_{\text{lo}} &= \frac{5}{4} \left. \left( \Pi^\mu \right) \right|_{\text{lo}} = \frac{5}{3} \left. \left( \Sigma^\mu \right) \right|_{\text{lo}}.
\end{align*}$$

Hence, the HTL contribution to $\Phi$ cannot be obtained by naively replacing the leading loop quantities in $\Phi^\mu$, as given above, by their HTL approximation. Instead, the HTL contribution $\Phi^\star$ can be conceived by analyzing how the $\mathcal{O}(e^2)$ discrepancy arises. Denoting the photon momentum in $\Phi$ by $K$ and the fermion momenta by $Q_{1,2}$, this diagram (with bare propagators for the $\mathcal{O}(e^2)$ contribution) can be represented as a double integral-sum over an expression with a numerator $N = K^2 - Q_1^2 - Q_2^2$. Closing the external legs of the boson self-energy in the HTL approximation amounts to neglecting the term $K^2$ in $N$. Tracing over the negative of the fermion HTL self-energy, on the other hand, neglects one of the $Q^2$ terms in $N$. Thus, all terms are accounted for twice in the sum over all three possibilities to approximate one of the momenta as soft. Accordingly, the representation

$$\Phi^\star = -\frac{1}{4} \text{Tr} D^\star \Pi^\star + \frac{1}{2} \text{Tr} S^\star \Sigma^\star$$

reproduces $\Phi^\mu$ perturbatively to order $\mathcal{O}(e^2)$. As shown in the following, this representation indeed leads to a well-defined resummed approximation of the thermodynamical potential.

\footnote{In the approach [3,4], instead, the functional properties of $\Phi$ in the leading loop approximation are used first to derive the entropy, which is then evaluated with the HTL propagators. It is noted, however, that in this framework the selfconsistency relation only holds approximately, see below.}
It is first emphasized that in the complete expression resulting from the HTL approximation of eq. (3) and (9), which can be written in a compact form as

$$\Omega^* = \frac{1}{2} \text{Tr} \left[ \ln(-D^* - 1) + \frac{1}{2} D^* \Pi^* \right] - \text{Tr} \left[ \ln(-S^* - 1) + \frac{1}{2} S^* \Sigma^* \right] - \Omega_{\text{ghost}},$$

with the individual contributions given by (7), (8), all temperature dependent divergences cancel. Moreover, the perturbative limit of the thermodynamical potential is reproduced by the representation (10). Separating the free contributions, e.g., for the boson part by

$$\ln(-D) = \ln(-D_0) + \ln(1 - D_0 \Pi),$$

and using the expansion

$$\ln(1-x) + x/2 = -x/2 + O(x^2),$$

the leading order correction to the interaction-free limit

$$\Omega_0 = -(d_T + \frac{7}{8} 2d_\pm) \frac{2^2}{90} V T^4$$

is

$$\Omega_0^{\pi} = -\frac{1}{4} \text{Tr} D_0 \Pi^* + \frac{1}{2} \text{Tr} S_0 \Sigma^* = \frac{T^2}{24} (M_b^2 + M_f^2) = \Omega^{pert}_{\pi}.$$ (11)

As in the HTL calculations \[6\] of the QCD entropy, the leading order term originates entirely from the behavior of the thermodynamically relevant excitations at the hard momentum scale $T$, which a posteriori justifies the present approach. However, in contrast to the entropy calculations \[3,4\] where the results are manifestly ultraviolet-finite, the cancellation of the thermal divergences is in the present approach directly related to the fact that the perturbative result is reproduced. It is emphasized that this aspect makes the representation (10) of $\Omega^*$ unique; any linear combination, apart from (9), of closed self-energy diagrams for the $\Phi$-contribution would result in either uncompensated thermal divergences or an incorrect perturbative limit.

The next-to-leading order term of the perturbative result $\Omega^{pert}$, on the other hand, cannot be expected to be reproduced in the present approximation: In an equivalent approach for the scalar $g^2 \phi^4$ theory, where the complete leading loop Luttinger-Ward resummation of the thermodynamical potential can be derived \[13\], the corresponding $O(g^3)$ correction takes its correct value only after the resummation of the self-energy, while the expressions (11) are calculated with bare propagators. Accordingly, the contribution of order $O(e^3)$, which arises from the static longitudinal parts of (10),

$$\Omega_{\text{nlo}} = \Omega_{A,h}^{\text{nlo}} + \frac{1}{2} \Omega_{B,h}^{\text{nlo}} = -\frac{(2M_b^2)^3/2 T}{12\pi} + \frac{(2M_f^2)^3/2 T}{16\pi}$$
is found to underestimate the perturbative result \( \Omega_{\text{nio}}^{\text{pert}} = -(e^2 T^2 / 3)^{3/2} T / (12 \pi) \) by a factor of 1/4. As observed in \([1,2]\) for QCD, the next-to-leading order term of \( \Omega_{A,b}^* \) agrees with the perturbative result but is overcompensated by the \( \Omega_{B,b}^* \) contribution, which is to be interpreted by a systematic next-to-leading order calculation.

V. SUMMARY

The generalized Luttinger-Ward representation of the thermodynamical potential is a suitable framework to derive consistently resummed approximations for relativistic gauge theories, with the propagators approximated by their HTL contributions. While the formalism is particularly simple to analyze in the case of the hot QED plasma, which is exemplified here, the application to other systems (including the QCD plasma) with the same HTL structure is evident. As a direct result of the HTL approximation, the resummed thermodynamical potential (10) is gauge independent. All medium-dependent divergences cancel, hence the approximation is explicitly renormalization-scale independent. The resummed expression (10) enjoys the anticipated behavior of a nonperturbative approximation. As shown in Figure 1, it yields a smooth extrapolation to the large-coupling regime, where it is inclosed by the perturbative results which are known to fluctuate with increasing order. In the weak-coupling limit, on the other hand, the perturbative result is recovered. This demonstrates that to leading order the resummed thermodynamical potential can entirely be expressed in terms of HTL propagators.

The applications of the formalism to the hot and dense quark-gluon plasma are straightforward and promising, in particular with regard to the suggestion \([14]\) to extrapolate finite-temperature lattice data to non-zero chemical potential. The reliability of the leading order HTL resummation of the thermodynamical potential in the large coupling regime, however, remains to be justified by a systematic next-to-leading order calculation.

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FIG. 1. The resummed pressure $p^* = -\Omega^*/V$ (solid line), as a function of the coupling and scaled by the interaction-free limit $p_0 = (2 + \frac{7}{8} \cdot 4) \frac{\pi^2}{90} T^4$, compared to the perturbative result in leading order (dashed line) and in next-to-leading order (dash-dotted line).
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