Medley in finite temperature field theory

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Abstract

I discuss three subjects in thermal field theory: why in $SU(N)$ gauge theories the $Z(N)$ symmetry is broken at high (instead of low) temperature, the possible singularity structure of gauge variant propagators, and the problem of how to compute the viscosity from the Kubo formula.

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I take this opportunity to make three comments about field theories at nonzero temperature, especially as applies to hot QCD. The first is trivial; the second, base speculation, and the third, a technical problem.

1. Why the $Z(N)$ symmetry is broken at high temperature

Consider an $SU(N)$ gauge theory without dynamical fermions. At nonzero temperature the order parameter for this system is the Polyakov line [1],

$$L(x) = \text{tr} \left( \mathcal{P} \exp \left( ig \int_0^\beta A_0(x) d\tau \right) \right);$$

$\mathcal{P}$ refers to path ordering, $g$ is the coupling constant, and $\beta = 1/T$, with $T$ the temperature. The Polyakov line is invariant under gauge transformations which are
strictly periodic in the euclidean time $\beta$. There are also gauge transformations which are periodic only up to an element of the center of the group $\mathbb{Z}$; the aperiodicity is a global $Z(N)$ phase, $\exp(i\phi_j), \phi_j = 2\pi j/N$ with $j = 0\ldots(N-1)$. The Polyakov line transforms under these aperiodic gauge transformations as $L(x) \rightarrow \exp(i\phi_j)L(x)$.

The Polyakov line has a nontrivial vacuum expectation value at high temperature, choosing one of $N$ equivalent vacua, and vanishes below the transition temperature. This behavior is exactly opposite to that of a typical spin system, where the symmetry is broken at low temperature, and restored at high temperature. Here I give a simple argument why.

In a spin system with discrete symmetry, the simplest argument for spontaneous symmetry breaking at low temperature is due to Peierls [3]. Consider an Ising model, with a spin $= \pm$ on each site, and a coupling $J$ between adjacent sites, so the hamiltonian $= -J \sum_{i,j} \sigma_i \cdot \sigma_j$. Start with a perfectly ordered state, where all spins $= +$. Now insert a domain of width $\sim R$ into this state, where all spins $= -$ inside the domain. Because the hamiltonian vanishes when the spins are aligned, the only energy is from the interface of this domain, $= JA$, with $A \sim R^{d-1}$ the area ($d$ is the number of spatial dimensions). This domain contributes to the partition function as

$$Z \simeq V e^{-JA/T}. \quad (1)$$

The factor of volume in the measure, $V \sim R^d$, is from the entropy of moving the domain around. Peierl’s argument is that in more than one dimension, at low temperatures the factor of $R^{d-1}$ in the exponential always wins over the $R^d$ in the measure: domains are exponentially suppressed, and the symmetry is spontaneously broken. In this view, it is not apparent how the symmetry is restored at a finite temperature; this happens because when domains are common they interact with one another, and this lowers their total energy [3].

$Z(N)$ interfaces can be constructed in hot gauge theories. In three spatial dimensions the interface tension between two $Z(N)$ domains is [4]

$$J \simeq \frac{4(N-1)\pi^2 T^3}{3\sqrt{3N}}. \quad (2)$$

The detailed origin of the interface tension is a story in itself. This result is computed semiclassically, and so is valid at high temperature, where the running coupling is small. The interface tension is the stationary point of an effective action, and represents a balance between classical and (certain) one loop effects. Because of this
balance, $J$ is proportional not to $1/g^2$, as is standard for instanton type solutions, but to $1/g$.

For my purposes here, however, the powers of $g$ are inconsequential — what matters is that the mass dimension is set entirely by powers of temperature, and not (say) by the renormalization mass scale. It is then obvious why the symmetry is broken at high instead of low temperature: the effective spin spin coupling is not constant (as for the Ising model), but increases with temperature, growing as $\sim T^3$ at large $T$. The contribution of a domain to the partition function is $Z \sim \exp(-\#T^2 R^2)$ at large $R$ and $T$: interfaces are rare at high temperature, and become common only as the temperature decreases.

What happens at large $N$, when $g^2 N$ is held fixed? The usual free energy of a gas of gluons is proportional to $N^2$, so it is natural to expect the same for the spin spin coupling $J$. Instead, one sees from (2) that if $g^2 N$ is held fixed at large $N$, that $J \sim N$ and not $N^2$. But remember that the $Z(N)$ vacua are for phases $\phi_j = 2\pi j/N$, so if two $Z(N)$ domains differ by a finite amount in $j$, the change in phase is only of order $1/N$. For two $Z(N)$ domains to differ by a finite phase, the difference in $j$ must be of order $N$. The total energy for such an interface is of order $N J$, which is of order $N^2$, as expected.

There is an amusing corollary to this picture. Consider the above argument at low temperatures. If the interface tension is nonzero as $T \to 0$, then Peierl’s argument implies symmetry breaking about zero temperature. Assuming that the symmetry is unbroken about zero temperature, then for the picture of $Z(N)$ domains to apply it is necessary for the interface tension to vanish at least linearly with temperature, $J \sim T$ as $T \to 0$.

I end this section with a matter of definition. I propose that the vacuum expectation value of the Wilson line is related to the free energy of a test quark, $F_q$, as

$$\langle L(x) \rangle = e^{-F_q/T+i\phi_j}.$$ (3)

The phase factor $\phi_j$ depends upon which $Z(N)$ domain the theory falls into. This relation is normally written without the phase factor (2). But then it is only consistent in the trivial phase, $\phi_j = 0$; otherwise, (3) implies that when $j \neq 0$, the free energy is complex, which is absurd.

The phase factor in (3) can be understood from the definition of the underlying partition function. To summarize: the Polyakov line represents not just the “mass” of the test particle (which at finite temperature is the free energy), but more gen-
erally, the propagator, \( L \sim \bar{\psi}(\beta)\psi(0) \). Now the gauge field, which resides in the adjoint representation, is invariant under the aperiodic gauge transformations mentioned previously. Fields in the fundamental representation, however, do change: while \( \psi(0) \rightarrow \psi(0) \), \( \bar{\psi}(\beta) \rightarrow \exp(i\phi_j) \bar{\psi}(\beta) \). This gives exactly the transformation expected for the Wilson line. More generally, one can show that the transformations of the Polyakov line under \( Z(N) \) transformations reflect the invariance of the states of the (purely gluonic) theory under global \( Z(N) \) gauge rotations [5].

2. Singularity structure of gauge variant propagators

In this section I conjecture what the possible singularity structure of gauge variant propagators are like in covariant gauges at nonzero temperature. The motivation is recent work by Baier, Kunstatter, and Schiff [3], who showed that special care must be taken in covariant gauges.

I first give a simple argument for the singularity structure of gauge variant propagators at zero temperature. Consider QED in covariant gauge, with gauge fixing parameter \( \xi \). Ignoring Dirac structure, at one loop order the fermion self energy for a fermion of mass \( m \) has the form

\[
\Sigma^{1\text{loop}}(\omega, k) \sim c e^2 (\omega - E_k) \int_{\omega - E_k}^m \frac{dk}{k} = (\omega - E_k) c e^2 \ln \left( \frac{m}{\omega - E_k} \right). \tag{4}
\]

The external energy is assumed to be near the mass shell, \( \omega \sim E_k \), and \( c = (3-\xi)/16\pi \). The logarithmic divergence is usually cut off by introducing a photon mass or some such. Instead, let me leave it as is, and assume that the effects of all higher order is simply to sum up the lowest order term into an exponential. Then the logarithmic divergence at one loop order becomes a branch point:

\[
\Delta^{-1}(\omega, k) = \omega - E_k + \Sigma^{1\text{loop}}(\omega, k) + \ldots \sim (\omega - E_k)^{1+ce^2}. \tag{5}
\]

While the above is just a guess, the result is correct: the correct singularity structure for the fermion propagator in QED is a branch point, where the strength of the singularity is gauge dependent [7].

At nonzero temperature I simply assume that the form is similar to that at zero temperature, with the addition of the appropriate factors for a thermal distribution:

\[
\Sigma^{1\text{loop}}(\omega, k) \sim (\omega - E_k) \xi \tilde{c} e^2 \int_{\omega - E_k}^T \frac{dk}{k} (1 + 2n(k)). \tag{6}
\]

Here \( n(k) = 1/(\exp(k/T) - 1) \) is the Bose-Einstein statistical distribution function, and \( \tilde{c} \) is a computable, nonzero constant. Note that at nonzero temperature the
gauge dependent term in (6) is strictly proportional to ξ, and vanishes if ξ does. This is because at nonzero temperature, the physical longitudinal and transverse modes acquire thermal “masses” which themselves cut off the infrared divergences near the mass shell. The sole exception to this are gauge dependent terms in covariant gauges; there the poles, massless at tree level, remain massless to any order in perturbation theory. Also, I assume the upper limit on any infrared divergence is now set by the temperature. It is then easy to see that for nonzero temperature, since the Bose–Einstein distribution function behaves as \( n(k) \sim T/k \) at small \( k \), that the integral for the self energy now has a linear divergence,

\[
\Sigma^{1\text{loop}}(\omega, k) \sim (\omega - E_k) \xi c e^2 \left( \frac{2T}{\omega - E_k} \right).
\]

This integral is precisely that studied by Baier, Kunstatter, and Schiff [6]: the factors of \( \omega - E_k \) appear to cancel, giving a gauge dependent contribution even on the mass shell. Baier, Kunstatter, and Schiff showed that this term has an imaginary as well as a real part, and so appears to give a gauge dependent contribution to the damping rate — which Braaten and I claimed was gauge invariant [8].

Rebhan was the first to point out that the above integral is infrared singular [9]: if computed with an infrared regulator, the term vanishes on the mass shell. This is simply because at one loop order there is a delicate balance between \( \omega - E_k \) upstairs against one over the same factor from the integral. Any infrared regulator will tend to smooth out the integral, so the \( \omega - E_k \) upstairs dominates, and \( \Sigma^{1\text{loop}} \) vanishes on the mass shell, \( \omega = E_k \).

My purpose here is merely to note that this is very special to one loop order and fails at higher order. Indeed, let us assume — as is true at zero temperature — that I can sum up the most infrared singular terms by exponentiation of the one loop result, to give

\[
\Delta^{-1}(\omega, k) \sim (\omega - E_k) \exp \left( 2\xi c e^2 T \frac{1}{\omega - E_k} \right).
\]

Because of the power like infrared singularities at nonzero temperature, instead of a branch point singularity there appears to be an essential singularity.

I do not have a grand moral from this calculation. Probably for reasons of prejudice as much as anything else, I believe that the position of the singularity in a gauge variant propagator is a gauge invariant quantity. Certainly the general proof of Kobes, Kunstatter, and Rebhan [10] does not depend on the detailed nature of the singularity in the inverse propagator: only that it vanishes on the mass shell. But the
tricks that work at one loop order probably will not be good enough beyond that. A more general methodology is needed.

3. The problem of viscosity

It’s really quite surprising that it’s so difficult to compute damping rates in hot gauge theories. Inevitably, this has bred a feeling of ennui and despair: after all, what are the damping rates good for? How do they affect measurable quantities?

One class of quantities which depend crucially on the damping rates are the transport coefficients [11,12]. In this section I repeat calculations of Ilyin et al. [11] in order to emphasize an apparent contradiction. Consider the Kubo formula for the shear viscosity, \( \eta \); the bulk viscosity, \( \zeta \), can be computed from \( \eta \). The Kubo formula for \( \eta \) is [11]

\[
\eta = \frac{1}{5} \lim_{E \to 0} \frac{1}{E} \text{Im} D(-iE + 0^+, 0). 
\]

\( D(p^0, \vec{p}) \) is the Fourier transform of a two point function of the stress energy tensor \( T^{ij} \): 

\[
D(\tau, \vec{x}) = <\pi^{ij}(\tau, \vec{x}) \pi^{ij}(0, 0)> ,
\]

where \( \pi^{ij} = T^{ij} - T^{kk} \delta^{ij}/3 \). The discontinuity is computed after analytic continuation from euclidean \( p^0 \) to \( p^0 = -iE + 0^+ \).

At lowest order the two point function of \( T^{ij} \) is given by the exchange of two transverse gluons,

\[
D(p^0, 0) = \frac{16}{3} \text{Tr} \Delta_t(k^0, k) \Delta_t(p^0-k^0, k) \left( 7k^4 - 10k^2k^0(p^0-k^0) + 7(k^0)^2(p^0-k^0)^2 \right). 
\]

(10)

Here \( k = |\vec{k}| \), \( \text{Tr} = T \sum_{k^0} \int d^3k/(2\pi)^3 \) is the integral over the loop momentum, and \( \Delta_t \) is the propagator for transverse gluons. The four powers of momenta arise because each \( T^{ij} \) brings in two derivatives of the gauge field. Since the stress energy tensor is gauge variant, gauge variant modes do not contribute to (10). The contribution of the plasmon modes is neglected. This is allowed because momenta of order \( T \) dominate, and in this range the plasmon modes are negligible (their residue is exponentially small).

Using the spectral representation of the transverse propagator,

\[
\Delta_t(k^0, k) = \int_0^{1/T} d\tau e^{ik^0\tau} \int^{+\infty}_{-\infty} d\omega \rho_t(\omega, k) (1 + n(\omega)) e^{-\omega\tau}, 
\]

(11)

where \( \rho_t(\omega, k) \) is the spectral density for transverse gluons, and \( n(\omega) \) the Bose–
Einstein statistical distribution function,
\[ \eta_{\text{gluon}}^{\text{loop}} = \frac{8}{15\pi T} \int_0^\infty k^2 dk \int_{-\infty}^{+\infty} d\omega n(\omega)(1 + n(\omega)) \rho_t^2(\omega, k) \left( 7k^4 - 10k^2\omega^2 + 7\omega^4 \right) \]  
(12)

The spectral density is assumed to be a Breit–Wigner form,
\[ \rho_t(\omega, k) = \frac{1}{2\pi k} \left( \frac{\gamma_t(k)}{(\omega - k)^2 + \gamma_t(k)^2} + \frac{\gamma_t(k)}{(\omega + k)^2 + \gamma_t(k)^2} \right), \]  
(13)

where \( \gamma_t(k) \) is the damping rate for a transverse gluon of momentum \( k \). Then
\[ \eta_{\text{gluon}}^{\text{loop}} = \frac{8}{15\pi^2 T} \int_0^\infty dk n(k)(1 + n(k)) \frac{k^4}{\gamma_t(k)} = \frac{32\pi^2 T^4}{225 \gamma_t}. \]  
(14)

The contribution of quarks can be computed similarly. For three flavors of massless quarks, the total result in QCD is
\[ \eta^{\text{loop}} = \frac{\pi^2}{225} \left( \frac{32}{\gamma_t} + \frac{63}{\gamma^+} \right). \]  
(15)

In this expression \( \gamma_t \) and \( \gamma^+ \) are the damping rates for gluons and quarks with momenta of order \( T \), which in that range are independent of momenta. Computation shows that these damping rates are of order \( g^2T \), so in all
\[ \eta^{\text{loop}} \sim \frac{T^3}{g^2}. \]  
(16)

The exact value for \( \eta^{\text{loop}} \) can be read off by using (15) and the damping rates in ref. [13]. What I am interested in here are the powers of \( g \).

At first sight it might seem odd that the shear viscosity diverges in the limit of a free theory, \( g \to 0 \). In fact this is a familiar feature of the shear viscosity: it is infinite for an ideal gas because if an ideal gas doesn’t start out precisely in thermal equilibrium, without interactions there is nothing to drive it there.

The problem of viscosity is the following. The stress viscosity can also be calculated using kinetic theory [12]. As is standard in kinetic theory, the result for \( \eta \) is one over the cross section. At lowest order the amplitude for the scattering of two gluons going into two gluons is of order \( g^2T \), and so
\[ \eta^{\text{transport}} \sim \frac{T^3}{g^4}. \]  
(17)

Much work has gone into computing the subleading logarithms in (17); Baym et al. [12] pointed out that the scale is set not by magnetic mass (as first thought) but by
Debye frequencies, of order $gT$. But forget the logarithms — even the powers of $g$ don’t match up between (16) and (17).

This discrepancy is special to hot gauge theories. Consider a scalar field theory with four point coupling $\lambda$. The damping rate first enters at two loop order, $\gamma \sim \lambda^2$, so there one expects that both Kubo and transport theory give $\eta \sim 1/\lambda^2$ [14].

It was pointed out to me first by Gordon Baym, and later by Eric Braaten, that $\eta_{\text{loop}}$ is incomplete. This is also discussed recently by Jeon [14]. The above calculation represents the resummation of an infinite set of diagrams, which are self energy corrections on two hard lines in a loop. Besides this diagram, there are also ladder diagrams, in which soft gluons are emitted, all parallel to one another, from the top to the bottom hard line in the loop. One can easily show by power counting that these diagrams are singular as $\gamma \to 0$, and so contribute to the stress viscosity.

If the kinetic theory calculation is correction, then the ladder diagrams must dominate, and give $\eta \sim 1/g^4$. Then $\eta_{\text{loop}}$ would be down by $g^2$. Using the Kubo formula would just be a dumb way of doing things.

Otherwise: the result is $\eta \sim 1/g^2$, with both ladder diagrams and $\eta_{\text{loop}}$ contributing at the same order; then there is more to the kinetic theory of ultrarelativistic plasmas than first meets the eye.

I suspect that the former is most probably true. But in either case, one can NOT conclude that the damping rates are irrelevant; they are certainly part of the answer for the transport coefficients.

Lebedev and Smilga [15] have studied the contribution of the damping rate to the photon self energy in hot QED. They find that in the sum of the hard loop, similar to $\eta_{\text{loop}}$, and ladder diagrams, that the damping rate cancels. While the damping rate cancels in the photon self energy, this does not imply that it does so in all quantities. Certainly there is a stress viscosity, computable (at least in principle) from the Kubo formula.

I conclude this section by showing that once one has the stress viscosity in hand, the bulk viscosity follows easily in a surprising way. The standard relation between the bulk and stress viscosities is

$$\zeta = \left(1 - 3 \frac{\partial p}{\partial \epsilon}\right)^2 \frac{5}{3} \eta. \tag{18}$$

Here $p$ is the pressure and $\epsilon$ the energy of the gas at a temperature $T$; calculation at next to leading order shows

$$p = T^4(a_1 - a_2g^2 + \ldots), \tag{19}$$
where \( a_1 \) and \( a_2 \) are constants which depend upon the number of colors and flavors \( [16] \). The energy is given from the pressure by \( \epsilon = T^2 \partial(p/T)/\partial T \). If the coupling constant is considered as a fixed parameter, then it is easy to see that even with the term \( a_2 \) in (19), that as for an ideal gas \( \partial p/\partial \epsilon = 1/3 \), and the bulk viscosity vanishes.

Of course the coupling is not a fixed parameter: it “runs”, and thereby develops a nontrivial dependence on temperature. With \( c_1 \) the lowest order coefficient of the \( \beta \)-function, \( \partial g/\partial \ln(T) = -c_1 g^3 \) \( (c_1 = (11 - 2N_f/3)/(16\pi^2)) \) for \( N_f \) massless flavors in \( SU(3) \) color,

\[
\zeta = \left( \frac{5a_2 c_1}{6a_1} g^4 \right) \frac{5}{3} \eta = \frac{273,375}{1,478,656} \frac{g^8}{\pi^8} \eta.
\]

The number refers to three massless flavors in \( SU(3) \) color.

The appearance of the lowest order coefficient of the \( \beta \)-function isn’t a complete surprise. For an ideal gas, or any system in which the trace of the stress energy tensor vanishes, \( \partial p/\partial \epsilon = 1/3 \). As is well known, for massless theories the conformal anomaly implies that the trace of the stress energy tensor is proportional to the \( \beta \)-function, which is why it enters into the relationship between the two viscosities.

For practical purposes the bulk viscosity is usually neglected relative to the stress viscosity. Certainly in the strict perturbative regime \( \zeta \) is negligible relative to \( \eta \). But the peculiar ratio of integers in front of \( \eta \) is about \( 1/5 \); thus \( \zeta \) is an extremely sensitive function of \( g \). As the coupling varies from \( g \sim 1 \) to \( g \sim 3 \), \( \zeta \) varies from essentially zero to a value commensurate with \( \eta \). Perhaps this enormous sensitivity has phenomenological consequences; after all, \( g \sim 3 \) is still \( \alpha = g^2/4\pi \) of order one.

In summary, while I cannot point to a direct experimental probe of the damping rates, they are crucial in sorting out the transport coefficients, and thereby deserving of our attention.

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