QCD PRESSURE AND THE TRACE ANOMALY

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Abstract

Exact relations between the QCD thermal pressure and the trace anomaly are derived. These are used, first, to prove the equivalence of the thermodynamic and the hydrodynamic pressure in equilibrium in the presence of the trace anomaly, closing a gap in previous arguments. Second, in the temporal axial gauge a formula is derived which expresses the thermal pressure in terms of a Dyson-resummed two-point function. This overcomes the infrared problems encountered in the conventional perturbation-theory approach.

1 Introduction

Conventional methods of calculating the QCD thermal pressure in perturbation theory have encountered an apparently insuperable infrared problem[1]. In previous papers, we have suggested a resummation technique that may overcome this problem and tested it on scalar field theory[2]. Although in principle this technique can be applied to QCD, in practice this would be far from simple, since it involves introducing a mass for the gluon. Although this mass must disappear at the end of the calculation, its presence at intermediate stages is a considerable inconvenience. In this paper, therefore, we introduce a more suitable technique for performing the resummation that is needed to overcome the infrared problem.

It is common to distinguish two different definitions of pressure[3]. The first, called the thermodynamic pressure, is given in terms of the grand partition function

$$ Z = \sum_i \langle i | \exp(-\beta H) | i \rangle $$

The summation is over a complete set of physical states $i$. The Hamiltonian $H$ is an integral over an element of the energy-momentum tensor:

$$ H = \int d^{n-1}x \, T^{00}(x) $$

The thermodynamic pressure is defined to be

$$ P = TV^{-1} \log Z $$
The other definition of the pressure is in terms of the thermal average of another element of the energy-momentum tensor. The hydrodynamic pressure is

\[ P = \langle T^{33} \rangle \]  \hspace{1cm} (1.4)

Although the two pressures \( P \) and \( \mathcal{P} \) are generally expected to be equal for a system in thermal equilibrium\[^{3}\], it was remarked some time ago by Landsman and van Weert\[^{4}\] that the presence of the trace anomaly might cause certain complications.

In section 2 we derive formulae for the derivatives of the thermodynamic pressure \( P \) with respect to the bare coupling \( g \) and to the temperature: see (2.7) and (2.8). In section 3 we replace \( g \) with the renormalised coupling \( g_R \), and we use the renormalisation group to show that the two formulae for the derivatives of the thermodynamic pressure \( P \) together imply that \( P \) is equal to the hydrodynamic pressure \( \mathcal{P} \). We show also that the pressure’s departure from simple \( T^4 \) behaviour is calculable from the QCD trace anomaly. In section 4 we show, by working in the temporal gauge, how simple Dyson resummation of the thermal gluon propagator removes the infrared problem. In section 5 we derive some simple results for the small-coupling limit, which will be the starting point for a perturbation expansion. We plan to consider such an expansion in a further paper.

2 Unrenormalised formulae

Consider pure-glue QCD for simplicity. At finite temperature there is a choice: one may heat up only the physical components of the gauge field\[^{5}\], but for our purposes it will be more convenient to adopt the more conventional formalism where also the unphysical components are heated. Then one may write\[^{6}\] the grand partition function as a path integral:

\[ Z(g, T) = \int dA d\bar{c} d\bar{c} dB \exp \left[ i \int_{\tau}^{\tau - i\beta} dx^0 \int d^{n-1}x \mathcal{L}(g, x) \right] \]  \hspace{1cm} (2.1)

The Lagrangian is

\[ \mathcal{L}(g, x) = -\frac{1}{4} F^2 + \mathcal{L}_{GF} + \mathcal{L}_{\text{GHOST}} \]  \hspace{1cm} (2.2)

We have chosen to implement the gauge fixing by means of an auxiliary field \( B \). The two common choices of \( \tau \) are \( \tau = 0 \), with the \( x^0 \) integration running along the imaginary axis, which is the imaginary-time formalism, and \( \tau = -\infty \), with the \( x^0 \) integration following the Keldysh contour, which is a version of the real-time formalism\[^{6}\]. Either formalism may be used for most of our work, though when a definite choice has to be made we will choose the second. So far, the fields are unrenormalised.

We scale each field by some power of \( g \), so obtaining new fields \( *A, \; *c, \; *\bar{c}, \; *B \). The integration measure then acquires a power of \( g \) that depends, loosely speaking, on the number of space-time points, so it is \( T \)-dependent. In order to cancel this, we consider the ratio

\[ \frac{Z(T, g)}{Z(T, 0)} \]  \hspace{1cm} (2.3)

and make the same field transformation in the denominator as in the numerator, so that the extra factor cancels. We then differentiate the logarithm of (2.3) with respect to \( g \), and transform back to the original fields. Choosing different powers of \( g \) in the definitions of the starred fields and using
the fact that the derivative of (2.3) must be independent of what powers we choose, we may obtain various identities. For our purposes, the most useful change of field variables is

$$A = A/g \quad c = c \quad \bar{c} = \bar{c} \quad B = Bg$$

(2.4)

We use the definition (1.3) of the thermodynamic pressure in terms of the grand partition function, and the translation invariance property that

$$\left\langle \int \tau^T \int d^{n-1}x F^2(x) \right\rangle = -i \beta V \left\langle F^2(0) \right\rangle$$

(2.5)

where $V$ is the volume of the $(n-1)$-dimensional $x$-space. We find that

$$\hat{P}(T, g) = P(T, g) - P_{\text{FREE}}(T)$$

(2.6)

satisfies

$$\frac{\partial}{\partial g} \hat{P}(T, g) = \frac{1}{2g} \left[ \left\langle F^2(0) \right\rangle - \left\langle F^2(0) \right\rangle_{\text{FREE}} \right]$$

(2.7)

Notice the importance of the subtraction term: it avoids a divergence at $g = 0$.

We may directly obtain a different derivative of $\hat{P}(T, g)$. We use the original definition (1.1) of $Z$, together with the expression (1.2) for the Hamiltonian $H$ as an integral over the energy density $T^{00}$:

$$\frac{\partial}{\partial \beta} \left[ \beta \hat{P}(T, g) \right] = - \left[ (T^{00}(0)) - \left\langle T^{00}(0) \right\rangle_{\text{FREE}} \right]$$

(2.8)

where again we have used translational invariance.

A gauge-invariant form for the gluonic part of the energy-momentum tensor is

$$T^{\mu\nu} = -F^{\mu\rho}F^{\nu}_{\rho} + \frac{1}{4}g^{\mu\nu}F^2$$

(2.9)

A trace over colour indices is understood. Classically, $T^{\mu\nu}$ is traceless, but in quantum theory there is an anomaly that changes this. We use dimensional regularisation so that we work initially $n = 4 - \epsilon$ dimensions. Then

$$T^\mu_{\mu} = -\frac{1}{4} \epsilon F^2$$

(2.10)

Renormalisation leaves behind a nontrivial limit as $\epsilon \to 0$.

Introduce the notation

$$D^M = \left\langle F^{M\rho}F^M_{\rho} \right\rangle - \left\langle F^{M\rho}F^M_{\rho} \right\rangle_{\text{FREE}}$$

(2.11)

where $M$ is not summed. Then

$$\frac{\partial}{\partial g} \hat{P}(T, g) = \frac{1}{2g} (D^0 - (3 - \epsilon)D^3)$$

(2.12a)

$$\frac{\partial}{\partial \beta} \left[ \beta \hat{P}(T, g) \right] = \frac{3}{4}D^0 + \frac{1}{4}(3 - \epsilon)D^3$$

(2.12b)

in $n = 4 - \epsilon$ dimensions.
3 Renormalisation

In (2.12), the coupling \( g \) is unrenormalised, and both \( D^0 \) and \( D^3 \) are divergent. We now replace \( g \) with a renormalised coupling \( g_R \). For this, we use dimensional regularisation and the MS scheme, so that we define the \( g_R \) by

\[
g = \mu^{\epsilon/2} Z(g_R) g_R(\mu)
\]

where \( Z(g_R) \) is a combination of wave-function and vertex renormalisation factors, and a beta function by

\[
\beta(g_R) = \mu \frac{\partial g_R}{\partial \mu} \bigg|_g
\]

The beta function has the structure\[7\]

\[
\beta(g_R) = -\frac{1}{2} \epsilon g_R + \tilde{\beta}(g_R)
\]

where for small \( g_R \)

\[
\tilde{\beta}(g_R) \sim \beta_0 g_R^3
\]

Differentiating (3.1a) with respect to \( \mu \) at fixed unrenormalised coupling \( g \), we obtain

\[
0 = \beta(g_R) \left[ \frac{\partial}{\partial g_R} \log Z(g_R) + \frac{1}{g_R} \right] + \frac{1}{2} \epsilon
\]

Differentiating it instead at fixed \( \mu \), we then have

\[
\frac{dg}{g} = \left[ \frac{\partial}{\partial g_R} \log Z(g_R) + \frac{1}{g_R} \right] dg_R = -\frac{\epsilon}{2\beta(g_R)} dg_R
\]

The pressure, being a physical quantity, is the same before and after renormalisation:

\[
P_R(T, g_R) = P(T, g)
\]

Of course, \( P_R \) depends also on \( \mu \). Hence (2.7) is equivalent to

\[
\beta(g_R) \frac{\partial}{\partial g_R} \hat{P}_R(T, g_R) = -\frac{1}{4} \epsilon \left[ \langle F^2(0) \rangle - \langle F^2(0) \rangle_{\text{FREE}} \right]
\]

\[
= \left[ \langle T^\mu_{\mu} \rangle > - < T^\mu_{\mu} \rangle_{\text{FREE}} \right]
\]

where we have used (2.10). So (2.12a) becomes

\[
\beta(g_R) \frac{\partial}{\partial g_R} \hat{P}_R(T, g_R) = -\frac{1}{4} \epsilon (D^0 - (3 - \epsilon)D^3)
\]

Because the left-hand side of (3.5) must be finite when \( \epsilon \to 0 \), the divergent part of \( (D^0 - 3D^3) \) as a function of \( g_R \) must be of order \( \epsilon^{-1} \) exactly. Because the derivative of the pressure in (2.12b) is a physical quantity and so must be finite, the divergent parts of \( D^0 \) and \( D^3 \), which are temperature-dependent, must be equal and opposite:

\[
D^0(T, g_R) \sim \frac{D(T, g_R)}{\epsilon} + d^0(T, g_R) + O(\epsilon)
\]

\[
D^3(T, g_R) \sim -\frac{D(T, g_R)}{\epsilon} + d^3(T, g_R) + O(\epsilon)
\]
Again, \(D^0\) and \(D^3\) depend also on \(\mu\). Hence when \(\epsilon \to 0\)
\[
\lim_{\epsilon \to 0} \frac{\partial}{\partial g_R} \hat{P}_R(T,g_R) = -\frac{1}{\beta(g_R)} D(T,g_R) = -\lim_{\epsilon \to 0} \frac{\epsilon}{\beta(g_R)} D^0
\]
(3.7a)
or, equivalently,
\[
\frac{\partial}{\partial g_R} \hat{P}_R(T,g_R) = -D(T,g_R) = \langle T^\mu \rangle
\]
(3.7b)

We have used the fact that \(< T^\mu >_{\text{FREE}}\) vanishes in the limit \(\epsilon \to 0\).

We now use the renormalisation group to show that the thermodynamic pressure \(P\) is equal to the hydrodynamic pressure \(\hat{P}\) defined in (1.4). Introduce the running coupling \(\bar{g}(T;g_R)\) such that
\[
T \frac{\partial}{\partial T} \bar{g}(T;g_R) \bigg|_{g} = \beta(g_R) \frac{\partial}{\partial g_R} \bar{g}(T;g_R) \bigg|_{T} \quad \bar{g}(\mu;g_R) = g_R
\]
(3.8a)

It is familiar that, to calculate \(\bar{g}(T;g_R)\), one makes use of the identity
\[
\beta(g_R) \frac{\partial}{\partial g_R} \bar{g}(T;g_R) \bigg|_{T} = \beta(\bar{g}(T;g_R))
\]
(3.8b)

Because \(\hat{P}_R\) has dimension \(T^n\), we may write
\[
\hat{P}_R(T,g_R) = T^n \phi(g_R,T/\mu)
\]
(3.9)

From renormalisation group theory,
\[
\phi(g_R,T/\mu) = \phi(\bar{g}(T;g_R),1)
\]
(3.10)

Thus
\[
T \frac{\partial}{\partial T} \phi(g_R,T/\mu) = \beta(\bar{g}(T;g_R)) \frac{\partial}{\partial \bar{g}(T;g_R)} \phi(\bar{g}(T;g_R),1)
\]
(3.11)

From (3.8b), we may replace \(\beta(\bar{g}(T;g_R))\partial/\partial \bar{g}(T;g_R)\) with \(\beta(g_R)\partial/\partial g_R\). Hence
\[
T^{n+1} \frac{\partial}{\partial T} \left[ T^{-n} \hat{P}_R(T,g_R) \right] = \beta(g_R) \frac{\partial}{\partial g_R} \hat{P}_R(T,g_R)
\]
(3.12)

With this equation, together with (2.8) and (3.5a), we obtain
\[
\hat{P}_R = -\frac{1}{n-1} \left\{ \left[ \langle T^\mu (0) \rangle - \langle T^\mu (0) \rangle_{\text{FREE}} \right] - \left[ \langle T^{00} (0) \rangle - \langle T^{00} (0) \rangle_{\text{FREE}} \right] \right\}
\]
\[= \frac{1}{n-1} \left[ \langle T^{ii} (0) \rangle - \langle T^{ii} (0) \rangle_{\text{FREE}} \right]
\]
\[= \left[ \langle T^{33} (0) \rangle - \langle T^{33} (0) \rangle_{\text{FREE}} \right]
\]
(3.13)

where in the last step we have used the spherical symmetry. From (1.4), this is just the statement that the thermodynamic pressure \(P\) is equal to the hydrodynamic pressure \(\hat{P}\). This result is valid in spite of the presence of the trace anomaly, so this closes a gap in the conventional proof that explicitly ignores it.[4]
This result is valid in spite of the presence of the trace anomaly. However, the anomaly does influence
the exact form of the pressure. In terms of the quantities introduced in (3.6), in the limit \( \epsilon \rightarrow 0 \)
\[
\hat{P} = -\frac{1}{4}(d^0 + d^3 - D)
\]  
(3.14)

The last term is just
\[
\frac{1}{4}D = -\frac{1}{4} \lim_{\epsilon \rightarrow 0} \langle T^\mu_\mu \rangle
\]  
(3.15)

We note that (3.12) and (3.5a) give also
\[
T \frac{\partial}{\partial T} \hat{P}_R(T, g_R) - (4 - \epsilon) \hat{P}_R(T, g_R) = \left[ \langle T^\mu_\mu(0) \rangle - \langle T^\mu_\mu(0) \rangle_{\text{FREE}} \right]
\]  
(3.16)

so when \( \epsilon \rightarrow 0 \) the departure of the pressure from simple \( T^4 \) behaviour is just the result of the QCD
trace anomaly. (The other parts of the energy-momentum tensor, arising from the gauge fixing and
the ghosts, do not contribute to this trace anomaly\([8]\).)

4 Temporal gauge

In the gauge \( A^0 = 0 \) the expression for \( D^0 \) in terms of the gluon field is particularly simple:
\[
D^0 = -\left[ \langle (\partial^0 A^i)^2 \rangle - \langle (\partial^0 A^i)^2 \rangle_{\text{FREE}} \right]
\]  
(4.1)

There are some unanswered questions about the temporal gauge, both at zero temperature\([9]\) and more
particularly in the finite-temperature imaginary-time formalism\([10]\), though the real-time formalism
may well be free of problems\([11]\). In the temporal gauge, we can ignore Faddeev-Popov ghosts.

With the Keldysh contour,
\[
D^0(0) = \int \frac{d^{4-\epsilon} q}{(2\pi)^{4-\epsilon}} \langle q^0 \rangle^2 [\Delta_{12}(q) - \Delta_{12,\text{FREE}}(q)]
\]  
(4.2)

where the subscript \( 12 \) refers to the element of the \( 2 \times 2 \) thermal matrix propagator\([6]\). This matrix
has the structure
\[
\Delta^{ij}(q) = M(q^0) \left\{ (\delta^{ij} - \frac{q^i q^j}{q^2}) \Delta_T(q) + \frac{q^i q^j}{q^2} \Delta_L(q) \right\} M(q^0)
\]  
(4.3)

where the matrices \( \Delta_T \) and \( \Delta_L \) are diagonal and
\[
M(q^0) = \sqrt{n(q^0)} \left[ e^{\frac{1}{\beta} |q^0|} \quad e^{-\frac{1}{\beta} q^0} \right] \quad \left[ e^{\frac{1}{\beta} q^0} \quad e^{\frac{1}{\beta} |q^0|} \right]
\]  
(4.4)

with \( n(q^0) \) the Bose distribution \( (e^{\beta |q^0|} - 1)^{-1} \). To zeroth order in the coupling
\[
\Delta_T(q) = \left[ \frac{\Delta_T(q)}{0} \right] \quad \Delta_T(q) = \frac{i}{q^2 + i\eta}
\]  
\[
\Delta_L(q) = \left[ \frac{\Delta_L(q)}{0} \right] \quad \Delta_L(q) = \frac{i}{(q^0)^2}
\]  
(4.5)
where some prescription is needed\textsuperscript{[9]} to define the meaning of $1/(q^0)^2$. As long as a rotationally invariant pole prescription is used, the Dyson resummed propagator has the same structure as in (4.3), but with

$$
\Delta_T(q) = \frac{i}{q^2 - \Pi_T(q, \beta)}
$$

$$
\Delta_L(q) = \frac{i}{(q^0)^2 - \Pi_L(q, \beta)}
$$

(4.6)

We need not retain the appropriate prescriptions $\Delta_L$ in (4.5); they play no part because the self-energies $\Pi_T$ and $\Pi_L$ have zero imaginary parts only at $q^0 = 0$, and when we use (4.6) in (4.2) the contribution from this point is killed by the $(q^0)^2$ in the numerator of the integrand. When we insert (4.2) into (3.7) we obtain

$$
\beta(g_R) \frac{\partial}{\partial g_R} \hat{P}_R(T, g_R)
= 8 \epsilon \operatorname{Im} \int \frac{d^4q}{(2\pi)^4} \theta(q^0)(q^0)^2 (1 + 2n(q^0)) \left\{ \frac{2}{q^2 - \Pi_T(q, \beta)} - \frac{2}{q^2} + \frac{1}{(q^0)^2 - \Pi_L(q, \beta)} - \frac{1}{(q^0)^2} \right\}
$$

(4.7)

The factor of 8 is included because a trace over colour indices was understood in all the equations. As in our previous work\textsuperscript{[2]}, any infrared divergence problem has been rendered harmless by the Dyson resummation: divergences of the self-energies are irrelevant because they now appear in denominators, and any zero of the denominators at $q = 0$ is unimportant because of the powers of $q$ present in $d^4q$.

In making use of (4.7) in a semi- or non-perturbative way, which is the goal of subsequent work, we have to renormalise accordingly. It is known\textsuperscript{[12]} that removing a factor $Z^{-2}$ from the first and third terms of the integral (4.7) renders their contribution to the integrand finite. This factor may be found by solving the differential equation (3.3a), with the boundary condition that $Z(0) = 1$. The solution may be written in the form

$$
\log Z(g_R) + \frac{1}{2} \log \left( \frac{2\beta(g_R)}{g_R \epsilon} \right) = 2 \int_0^{g_R} d\gamma \frac{\tilde{\beta}(\gamma)/\gamma^2 - \frac{1}{2}(\partial/\partial \gamma)(\tilde{\beta}(\gamma)/\gamma)}{\epsilon - 2\beta(\gamma)/\gamma}
$$

(4.8)

This integral converges when $\epsilon \to 0$, so the divergent part of $Z^{-2}$ is

$$
\frac{1}{Z^2} \sim - \frac{2\beta(g_R)}{g_R \epsilon}
$$

(4.9)

So (4.7) finally becomes

$$
\frac{\partial}{\partial g_R} \hat{P}_R(T, g_R) = -16 \operatorname{Lim}_{\epsilon \to 0} \epsilon \operatorname{Im} \int \frac{d^4q}{(2\pi)^4} \theta(q^0)(q^0)^2 (1 + 2n(q^0)) \left\{ \frac{2}{q^2 - \Pi_T^C(q, \beta)} - \frac{2}{q^2} + \frac{1}{(q^0)^2 - \Pi_L^C(q, \beta)} - \frac{1}{(q^0)^2} \right\}
$$

(4.10)

where $\Pi_T^C$ and $\Pi_L^C$ are convergent thermal self-energies.
5 Small-coupling limit

So far, our formulae are exact: we have not used perturbation theory to derive them. We now investigate how they behave in the limit of small coupling. We content ourselves with simple results that follow without detailed calculation, and leave a more extensive investigation to a future paper.

Because\[^{[1][13]}\] $P_R(T, g_R)$ behaves as $-\frac{1}{6} g_R^2 T^4$ for small $g_R$, we see from (3.2b) and (3.7b) that

$$D = O(g_R^4) \quad (5.1)$$

We may obtain almost the same information directly from (2.12a), which becomes with (3.6)

$$\frac{\partial}{\partial g} P_R(T, g) = \frac{1}{2g} \left[ \frac{4D}{\epsilon} \epsilon + d^0 - 3d^3 - D + O(\epsilon) \right] \quad (5.2)$$

To lowest order the unrenormalised and renormalised couplings are equal, so to this order the left-hand side is again $\partial/\partial g R P_R(T, g_R)$; therefore the singular term on the right-hand side must be at least of order $g_R^3$. Furthermore, in order to give agreement with (3.7b),

$$d^0 - 3d^3 = -\frac{2g_R D}{\tilde{\beta}(g_R)} + \ldots \quad (5.3)$$

where the further terms are at least of order $g_R^3$.

The pressure may also be calculated directly from (3.14). For this to be compatible with (5.2) in lowest order, it must be that $d^0 = d^3$ to lowest order. That is

$$P_R = -\frac{1}{2} d^0 + \ldots \quad (5.4)$$

where the further terms are of order $g_R^3$ at least.

6 Conclusion

(3.7) is an exact relation between the QCD thermal pressure and the thermal average of the trace anomaly. We have used this to prove nonperturbatively the equivalence of the thermodynamic and the hydrodynamic pressure.

With the plausible assumption that thermal field theory may be formulated consistently in the temporal gauge, (3.7) can be expressed in terms of a manifestly infrared finite integral (4.10) over the full propagator. The absence of a chromomagnetic screening mass poses no particular problem to an evaluation of (4.10), so it may provide a suitable starting point for developing new resummation techniques.

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