The Thomas-Fermi approximation for gauge theories

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ABSTRACT

An effective field approximation, similar to the atomic Thomas-Fermi approach, is proposed for studying non-Abelian gauge theories which includes finite-volume effects. As applications of the formalism the equation of state for an SU(2) gauge theory with massless fermions is obtained. The extensions to realistic situations are briefly discussed.

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1. Introduction

The equation of state of a quarks gluon plasma at high temperatures and/or densities is one of the most important unknowns in our current understanding of strong interaction physics\(^1\). The applications of such an equation of state are varied, ranging from cosmological compact objects to the physics of heavy ion collisions. Unfortunately, due to the high degree of non-linearity present in QCD, the determination of this equation of state has proved to be a difficult task. For example, even within perturbation theory, infrared singularities require the calculation of an infinite number of graphs for the partition function beyond fifth order\(^2\). General expressions for Green’s functions are available for the case where the internal momentum is large (\(\sim T\)) while the external momenta are soft (\(\sim gT\)), the so-called hard-thermal loop region\(^3\); using standard manipulations\(^4\), one can then determine the partition function corresponding to all hard modes in the theory. The soft-mode contributions to the partition function have been studied using various approximations\(^5\) and numerical calculations have also been developed (though not to the extent as in the \(T = 0\) case)\(^6\).

In this paper we propose a new approximation within which the physics of a quark-gluon plasma can be studied. The formalism is based on the Thomas-Fermi model of the atom\(^7\) and will be called Thomas-Fermi QCD (TFQCD). We consider a plasma of quarks and gluons confined to a volume \(V\) which we imagine subdivided into a number of subvolumes, each of which is large enough for the partons they contain to be considered a statistical ensemble. These subvolumes interact via a background gauge field whose sources are the thermally-averaged non-Abelian charge densities of the subvolumes. The subvolumes are assumed
to be small enough for the background field to vary very little inside them, and
because of this the background field sources are essentially point-like. The require-
ment of stability, together with the Yang-Mills equations for the background field,
furnishes a closed set of equations which can be solved; from the solution the equa-
tion of state for the system can be derived. This program requires the evaluation
of the thermally-averaged non-Abelian charge densities which we obtain using per-
turbation theory; in this paper we will use the lowest-order approximation, but a
systematic improvement is straightforward.

The atomic Thomas-Fermi approximation is useful when calculating bulk prop-
erties of an atom with a large number of electrons, such as the total ionization
energy\(^8\); it is also useful as a starting point for a Hartree-Fock approximation. We
expect the TFQCD model of a quark-gluon plasma to be reasonably accurate for
bulk properties of the system, such as the equation of state. There are some differ-
ences between the atomic Thomas-Fermi and the TFQCD formalisms; in particular
note that, in contrast to the atomic case, the quark-gluon plasma is not stable: if
left alone it will fly apart and undergo a phase transition into a gas of hadrons. In
order to study a gas of quarks and gluons we are forced to imagine the system to
be enclosed in a container at sufficiently large temperature and/or density.

The presence of an external confining agency is reminiscent of the bag model\(^9,10\).
Through most of the paper we will consider, in contrast to the usual bag models, a
situation where the partons are not confined, and for which the external pressure
is assumed to be generated by a some physical apparatus. Despite this difference
the bag boundary conditions are also relevant for the present model: the system is
assumed to be confined to a spherical volume out of which neither fermion number
nor color can escape, this requires we impose both the original\textsuperscript{9} and chiral\textsuperscript{10} bag boundary conditions.

We will also briefly study a system corresponding to a hadron at zero temperature, and will show that the bag constant and strong coupling constant obtained in the present approach are consistent with those obtained using the bag model.

The volume of the system $V$ will be kept finite in all computations; the results will then include finite-volume effects (such as terms in the extensive thermodynamic quantities proportional $V^{2/3}$). In the infinite volume limit these surface effects can be neglected and the equation of state reduces to that of an ideal gas of gluons and quarks.

In the following section we will describe the construction of the TFQCD model and present some simple applications. We will concentrate on the case of an $SU(2)$ gauge theory with a single species of massless fermions, and then describe the modification required for the important case of an $SU(3)$ gauge theory with three (massive) fermion flavors. The remainder of the paper is organized as follows. In section 3 we derive the equation of state for this model in the cases of zero baryon number and zero temperature. The discussion of the extension to $SU(3)$ and to more flavors is presented in section 4. Some parting comments are presented in section 5; and a mathematical detail is given in the appendix.
2. Description of the model

The model we propose is, as mentioned above, an extension of the Thomas-Fermi model to the case of QCD. We consider a gas of partons inside a volume \( V \); we then imagine partitioning \( V \) into small subvolumes \( \delta V \) which are big enough so that the partons (quarks and gluons) contained in them form a statistical ensemble determined by a temperature \( T \) and, for the fermions, a chemical potential \( \mu \). Each subvolume is required to be in equilibrium with its environment which implies that the temperature and chemical potential are the same throughout the system (this is intuitively obvious, we present a proof in Sect. 2.3). The system is also assumed to be static so that no currents are present.

We assume that the subvolumes have a non-zero average color charge, which implies that the zero-component of the gauge field goes to a constant, \( \bar{A}^0 \) at its boundary\(^{11} \). We will refer to \( \bar{A}^0 \) as the background gauge field. The background field is assumed to vary slowly and smoothly between the \( \delta V \), and is determined self-consistently by requiring it to satisfy the Yang-Mills equations corresponding to the average charges of the subvolumes (which themselves depend on the background fields). This approach presupposes that the magnitude of the charge in any given \( \delta V \) is small, and that the background field is approximately constant within each subvolume; both these assumptions will be verified \textit{a-posteriori}.

Finally, we also assume that our system is spherically symmetric; this requirement considerably simplifies the calculations yet preserves the essential non-Abelian character of the problem. The equations obtained for the background fields are then similar to the ones derived when considering the coupling of classical,
spherically symmetric Yang-Mills fields to external sources\textsuperscript{12,13}.

In the rest of this section we will treat the various ingredients of the model separately. We first review the Yang-Mills equations within the spherically symmetric Ansatz. We then obtain the expression for the partonic sources for the background fields and the various thermodynamic observables. Next we derive the stability conditions for the system. Finally we combine these results in order to obtain the equations for the background fields which determine quantitatively the Thomas-Fermi-QCD (TFQCD) model.

The conventions which we use are the following. The model is based on an $SU(N)$ Yang-Mills theory with one species of massless fermion; the (anti-hermitian) group generators are denoted by $T^a$ and the gauge coupling constant by $g$. The covariant derivative is $D_\mu = \partial_\mu + A_\mu$ where $A_\mu = g A^a_\mu T^a$. The full Lagrangian is

$$\mathcal{L} = i \bar{\psi} \slashed{D} \psi - \frac{1}{4} (F^a_{\mu\nu})^2$$  \hfil (2.1)

where $\psi$ denotes the quark field, $A$ the gauge field, and

$$F^a_{\mu\nu} = \partial_\mu A^a_\nu - \partial_\nu A^a_\mu + g \epsilon_{abc} A^b_\mu A^c_\nu.$$  \hfil (2.2)

The sources are

$$j^a_\mu = i \bar{\psi} T^a \gamma_\mu \psi.$$  \hfil (2.3)

Latin indices from the beginning of the alphabet ($a, b, c$, etc) correspond to color indices; Latin indices form the middle of the alphabet ($i, j, k$, etc) denote space indices.
2.1. Spherically symmetric gauge potentials and equations

As mentioned in the previous section we will assume that the long-range forces in our system are described by a non-Abelian background gauge field generated by the average charge of each subvolume. We also assume the system to have spherical symmetry. Thus we need the most general expression for a spherically symmetric non-Abelian gauge field, which is well known\textsuperscript{14}, and is reviewed for completeness below.

The most general spherically symmetric Ansatz for the gauge potentials of an Abelian theory is simply $A^0 = \phi(r, t), \ A = a(r, t) \ \hat{r}$, where $\hat{A}$ denotes the vector potential and $r = |r|$. It is clear, however, that we can choose a gauge where $a(r, t) = 0$, so we can take $\hat{A} = 0$.

For the $SU(2)$ non-Abelian case the structure is much richer\textsuperscript{#1}; the most general spherically symmetric Ansatz is\textsuperscript{14} (the over-bar denotes the background fields)

\[
\bar{A}^0_a = A_0 \hat{r}_a \\
\bar{A}^i_a = \epsilon_{iaj} \hat{r}_j \left( \frac{g^{-1} + \varphi_2}{r} \right) + (\delta_{ia} - \hat{r}_i \hat{r}_a) \frac{\varphi_1}{r} + \hat{r}_i \hat{r}_a A_1
\]

which exhibits spin-isospin mixing\textsuperscript{#2}. The fields $\varphi_{1,2}$ and $A_{0,1}$ depend on $r$ and $t$.

Within this Ansatz the $SU(2)$ Yang-Mills Lagrangian becomes

\[
\frac{1}{4} F^{a \mu \nu} F_{a \mu \nu} = \frac{1}{4} f^{2 \mu \nu} - \frac{1}{r^2} |D \Phi|^2 + \frac{g}{2r^4} \left( |\Phi|^2 - \frac{1}{g^2} \right)^2
\]

where the indices $\mu, \nu$, etc. equal 1 (corresponding to $r$) or 0 (corresponding to $t$);

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\textsuperscript{#1} The situation is similar for larger groups, see section 4.
\textsuperscript{#2} In this respect the present approach differs from other investigations into spherically symmetric hadron physics; see Ref. 15.
the metric is diag(1, −1). We also defined $f_{\mu\nu} = \partial_{\mu}A_{\nu} - \partial_{\nu}A_{\mu}$, $\Phi = \varphi_1 + i\varphi_2$ and $D_{\mu} = \partial_{\mu} + igA_{\mu}$. The above expression is invariant under the gauge transformation

$$A_{\mu} \rightarrow A_{\mu} - \partial_{\mu}\Lambda, \quad \Phi \rightarrow e^{ig\Lambda}\Phi,$$

which is a remnant of the original non-Abelian invariance.

We now consider the coupling of the above fields to a spherically symmetric charge density $\rho^a$, where spherical symmetry requires $\rho^a = q\tilde{r}^a$. The coupling is then described by adding a term

$$L_{\text{interaction}} = g\tilde{A}_0^a \rho^a = g q A_0$$

(2.7)

to the Lagrangian.

2.2. THE PARTITION FUNCTION FOR GLUONS AND FERMIONS

We now imagine that the volume of the system, denoted by $\mathcal{V}$, is subdivided into a large number of subvolumes $\delta\mathcal{V}$. The gauge fields inside each subvolume are separated into a background piece $\tilde{A}_\mu^a$ and a fluctuation $a_{\mu}^a$: $A_\mu^a = \tilde{A}_\mu^a + a_{\mu}^a$.

In this subsection we evaluate the partition function for the partons inside $\delta\mathcal{V}$. This object, which we call $Z_{\delta\mathcal{V}}$, will depend on $A_\mu^a$, and we can use this dependence to obtain the thermal average of the non-Abelian currents,

$$\tilde{\bar{j}}_\mu^a = \frac{1}{g} \left( \frac{\partial Z_{\delta\mathcal{V}}}{\partial A_\mu^a} \right)$$

(2.8)

Since the system is supposed to be in a static configuration we require $\tilde{\bar{j}}_i^a = 0$ which implies we can take $\tilde{A}_i^a = 0$ inside $\delta\mathcal{V}$. Since the background fields are
assumed slowly varying, we also take $\bar{A}_0^a$ constant inside $\delta V$; we then choose a gauge such that $\bar{A}_0^a$ is diagonal inside $\delta V$. Hence $Z_{\delta V}$ will depend on the temperature $T$, the fermionic chemical potential $\mu$, and the $n$ components of $\bar{A}_0$ associated with the diagonal generators ($n$ is the rank of the gauge group).

As a first approximation we will neglect the interaction between the fermions and the $a^a_\mu$, as well as the non-linear couplings among the $a^a_\mu$; these interactions can be included perturbatively.

Concerning the scale of $\delta V$ we will assume that it is set by the fermion thermal wavelength, $\lambda$, that is,

$$\delta V \sim \lambda^3. \quad (2.9)$$

We will for the moment restrict ourselves to the case where the gauge group is $SU(2)$ (the extension to $SU(3)$ will be described in section 4 below). In this case the group generators are $T^a = \sigma^a/(2i)$ and, within $\delta V$, $\bar{A}_0 = g\bar{A}_0^a T^a = gA_0 \sigma_3/(2i)$, so that

$$A_0^b = a_0^b + A_0 \delta_{b3}, \quad A_i^b = a_i^b \quad \text{(inside } \delta V) \quad (2.10)$$

We first evaluate the fermionic contribution to the partition function, and then calculate the contributions from the $a^a_\mu$.

**Fermionic contribution.**

When considering the fermionic partition function we will assume only one massless fermions species (the modifications required by several species and/or non-zero masses are straightforward). Thus we look for an approximate expression
for $Z_\psi = \det (i \slashed{D} + \mu \gamma_0)$ where the gluon fields take the form (2.10). Inside a subvolume $\delta V$ it is assumed that the fermions behave as a statistical ensemble, that the interaction with the $a^a_\mu$ is small, and that the background fields $\bar{A}^a_\mu$ are essentially constant. Adopting these approximations we reduce the calculation to evaluating $\det \left[ i \slashed{D} + \left( -i \bar{A}_0 + \mu \right) \gamma_0 \right]$ with $\bar{A}_0 = gA_0 \sigma_3 / (2i)$, $A_0 = \text{constant}$.

The partition function for an ideal gas of massless fermions at temperature $T$ in a volume $\delta V$ and with a chemical potential equal to $\mu$ is given by:

$$\ln Z_0 = \frac{\beta \delta V}{12\pi^2} \left[ \mu^4 + 2(\pi kT)^2 \mu^2 + \frac{7}{15}(\pi kT)^4 \right]$$

(2.11)

where the zero subscript indicates that no gauge fields are included. The constant background gauge fields are then included by replacing $\mu \rightarrow \mu \pm gA_0/2$ in $Z_0$ with the sign depending on the isospin of fermion, and adding the contributions from each isospin component. Thus, within the above approximations, we obtain

$$\ln Z_\psi = \frac{\beta \delta V}{6\pi^2} \left[ \frac{g^4 A_0^4}{16} + \frac{g^2 A_0^2}{8\lambda^2} + \mu^4 + 2 \left( \pi kT \mu \right)^2 + \frac{7}{15} \left( \pi kT \right)^4 \right]$$

(2.12)

where we defined the thermal wavelength

$$\lambda = \frac{1}{2 \sqrt{3\mu^2 + (\pi kT)^2}}$$

(2.13)

This approximation to $Z_\psi$ generates the following expressions for the fermionic
contribution to the (local) thermodynamic quantities

\[ P_\psi = \frac{\ln Z_\psi}{\beta \delta V} = \frac{1}{6\pi^2} \left[ \frac{g^4 A_0^4}{16} + \frac{g^2 A_0^2}{8\lambda^2} + \mu^4 + 2(\pi kT\mu)^2 + \frac{7}{15}(\pi kT)^4 \right], \]

\[ s_\psi = \frac{2k}{3\pi}(\pi kT) \left[ \frac{g^2 A_0^2}{4} + \mu^2 + \frac{7}{15}(\pi kT)^2 \right], \]

\[ e_\psi = \frac{1}{2\pi^2} \left[ \frac{g^2 A_0^2}{24\lambda^2} - \frac{g^4 A_0^4}{48} + \mu^4 + 2(\pi kT\mu)^2 + \frac{7}{15}(\pi kT)^4 \right], \]

\[ n = \frac{1}{\beta \delta V} \left( \frac{\partial \ln Z_\psi}{\partial \mu} \right) = \frac{2\mu}{3\pi^2} \left[ \frac{3}{4} g^2 A_0^2 + \mu^2 + (\pi kT)^2 \right], \]

\[ q_\psi = \frac{1}{g \beta \delta V} \left( \frac{\partial \ln Z_\psi}{\partial A_0} \right) = \frac{g A_0}{24\pi^2} \left[ g^2 A_0^2 + \frac{1}{\lambda^2} \right], \]

where \( P \) denotes the pressure and \( s, e, n, q \), etc denote the entropy, energy, particle and charge per unit volume. Note that \( e_\psi = 3P_\psi - 2A_0q \) includes the energy of the interaction with the gauge fields.

Gluonic contribution

The gluonic contribution to the partition function is obtained in a manner similar to the one followed for the fermions. As before we will ignore the self-interactions of the fields \( a_{\mu}^{a} \), in this case the partition function, including the Fadeev-Popov determinant, reduces to \( Z_{\text{gluons}} = \text{det} - D_{\text{adj}}^2 \) where \( D_{\text{adj}} \) denote the covariant derivative for the background fields in the adjoint representation\(^{16}\).

In calculating this determinant we take into account the partons are supposed to be in a box of side \( \sim \lambda \). Moreover, the background field is supposed to include the effects from the zero (Fourier) modes in the field. It follows that we need to include only modes with energy above

\[ p_o = \frac{2\pi}{\lambda}, \]

(2.15)
Using the gauge $A_\mu = \delta_{\mu,0} \delta^{a,3} A_0$ gives

$$\ln Z_{\text{gluons}} = -V \int_{p > p_o} \frac{d^3 p}{(2\pi)^3} \left[ \ln \left( 1 - e^{-\beta(p_g A_0)} \right) + \ln \left( 1 - e^{-\beta(p + g A_0)} \right) + \ln \left( 1 - e^{-\beta p} \right) \right]$$

(2.16)

which corresponds to a gas of massless bosons with chemical potential $\pm g A_0$ and 0.

We will argue below (section 2.4.1) that the background field $A_0$ is monotonic in $r$ and that $\lambda g A_0 \leq 3\pi/2$. Using also the fact that $\beta p_o \geq 4\pi^2$ we find that to a good approximation

$$P_{\text{gluons}} = \frac{(\pi kT)^4}{15\pi^2} + \frac{4 (kT)^4}{\pi^2} - \left( \beta^2 p_o^2 + 2 \beta p_o + 2 \right) e^{-\beta p_o} \left[ \sinh \left( \beta g A_0 / 2 \right) \right]^2$$

(2.17)

This shows that the deviations from the free-gluon values are exponentially suppressed (recall that $\beta p_o \geq 4\pi^2$) and can be neglected. In this case

$$P_{\text{gluons}} \approx \frac{(\pi kT)^4}{15\pi^2}, \quad e_{\text{gluons}} \approx \frac{(\pi kT)^4}{5\pi^2}, \quad s_{\text{gluons}} \approx \frac{4k(\pi kT)^3}{15\pi}, \quad q_{\text{gluons}} \approx 0$$

(2.18)

where $P$ denotes the pressure, and $e$, $s$ and $q$ the energy, entropy and charge per unit volume respectively. The errors incurred are below a few percent for the thermodynamic quantities and below 0.0075% for the charge.

2.3. Stability conditions

The stability criterion can be obtained from the Wong equations\textsuperscript{17}, but a more elegant argument can be gleaned from a paper by Brown and Weisberger\textsuperscript{18}.
Consider the background field contribution to the energy momentum tensor $\theta_{\text{back}}$, which satisfies
\[ \partial_{\mu} \theta_{\text{back}}^{\mu\nu} = g \rho_a \bar{F}_a^{\nu\mu} . \] (2.19)
where $\bar{F}_a^{\nu\mu}$ denotes the field strength for the background fields, and $\rho_a$ the thermally-averaged non-Abelian charge density.

Since the total energy momentum tensor is conserved, it follows that the averaged partonic contribution $\theta_{\text{part}}$ satisfies
\[ \partial_{\mu} \theta_{\text{part}}^{\mu\nu} = -g \rho_a \bar{F}_a^{\nu\mu} . \] (2.20)
For static situations the above equation implies
\[ \partial_i \theta_{\text{part}}^{ij} = -g \rho_a \bar{F}_a^{ij} . \] (2.21)
If in addition we impose spherical symmetry (see section 2.1) $\rho_a \propto \hat{r}^a$, $\bar{A}_a^0 \propto \hat{r}^a$ which implies $\rho_a \bar{F}_a^{ij} = -\rho_a \partial_j \bar{A}_a^0$. For a homogeneous gas of partons the space components of the energy momentum tensor are $\theta_{\text{part}}^{ij} = P_{\text{part}} \delta_{ij}$. Collecting these results we get $\partial_j P_{\text{part}} = g \rho_a \partial_j \bar{A}_a^0$, or, equivalently
\[ dP_{\text{part}} = g \rho_a d\bar{A}_a^0 \] (2.22)
which is the desired constraint.

For the Abelian case (2.22) reduces to the usual Thomas-Fermi equilibrium condition: the pressure on $\delta \mathcal{V}$ is balanced by the electrostatic force.
This stability condition requires the chemical potential and temperature to be \( r \) independent. Indeed, \( \ln Z_{\delta V} \), the parton partition function for a small volume \( \delta V \), is a function of \( T, A_0 \) and \( \mu \); using (2.14) we obtain

\[
dP_{\text{part}} = \frac{1}{\beta V} \left( \frac{\partial \ln Z_{\delta V}}{\partial A_0} \right) dA_0 + \frac{1}{\beta V} \left( \frac{\partial \ln Z_{\delta V}}{\partial \mu} \right) d\mu + \frac{1}{\beta V} \left( \frac{\partial \ln Z_{\delta V}}{\partial T} \right) dT
\]

\[
= g q dA_0 + n_{\text{part}} d\mu + (e_{\text{part}} - \mu n_{\text{part}}) \frac{dT}{T}
\]

(2.23)

where \( n_{\text{part}} \) and \( e_{\text{part}} \) are, respectively, the particle and energy densities of the partons, and \( \rho^a = q \hat{r}^a \). Substituting this expression for \( \rho_a \), using (2.4), and comparing to (2.22) we obtain \( d\mu = dT = 0 \).

2.4. The TFQCD equations

The equations of motion are derived from the spherically symmetric Lagrangian for the background fields (2.5) when the potentials interact with a source \( q \) according (2.7). The source, given in (2.14), is itself a function of the potentials. The resulting equations are

\[
D^2 \Phi + \frac{g^2}{r^2} \left( |\Phi|^2 - \frac{1}{g^2} \right) \Phi = 0
\]

\[
\partial^\mu \left( r^2 f_{\mu\nu} \right) + 2g \text{Im} (\Phi^* D_\nu \Phi) = -g r^2 q_\nu^0.
\]

(2.24)

The gauge invariance of these equations allows us to chose the \( A_1 = 0 \) gauge. The second of the above equations gives, when \( \nu = 1 \) and for static configurations, \( \text{Im} \Phi^* \Phi' = 0 \), so that we can choose \( \Phi \) to be purely imaginary.

We will use the notation

\[
A_0 = \frac{f(r)}{rg}, \quad \Phi = \frac{1}{ig} a(r).
\]

(2.25)
Then, using (2.14), the above equations become

\[ f'' - 2 \left( \frac{a}{x} \right)^2 f = \frac{\alpha}{6\pi} f \left( \frac{f^2}{x^2} + 1 \right) \]

\[ a'' + \frac{1 + f^2 - a^2}{x^2} a = 0. \tag{2.26} \]

where \( \alpha = g^2 / (4\pi) \) and \( x = r/\lambda \) with \( \lambda \) defined in (2.13).

These equations determine the background self-consistently. Their solution requires the specifications of the boundary conditions to which we now turn.

**Boundary conditions** The conditions near the origin are determined by considering the behavior of Wilson loops as \( r \to 0 \). We find that singularities arise unless \( f \) and \( a^2 - 1 \) vanish at \( r = 0 \). Using the \( a \leftrightarrow -a \) symmetry (which is a remnant of the gauge symmetry) we can then require \( a \to 1 \) as \( r \to 0 \). The precise manner in which \( f \) and \( a - 1 \) vanish as \( r \) approaches zero is determined by requiring that the energy should have no divergences at this point. We then obtain

\[ f, \quad a - 1 = O(r^2) \quad \text{for} \quad r \to 0. \tag{2.27} \]

It is easy to see that the equations of motion (2.26) require \( f \) to be concave or convex; since we can always exchange \( f \to -f \) we can assume that \( f \) is concave. In particular this implies that \( f \) will not vanish for \( r > 0 \). From (2.26) it is also easy to show that \( (f/r)' > 0 \). In contrast \( a \) can (and will) have extrema as well as zeroes.

As mentioned in the introduction we assume that the system is enclosed in a container which must be spherical due to requirement of spherical symmetry; we
denote by \( R \) its radius. If the system is to be confined to the region \( r < R \), there should be no leakage of fermion number or color into the region \( r > R \).

The first of these two conditions (fermion number confinement) requires the fermions to have zero radial component of the momentum at \( r = R \). This implies that in the vicinity of \( r = R \) the fermion gas becomes two dimensional. The corresponding (surface) charge density \( \sigma^a \) takes the form

\[
\sigma^a = \vartheta \tilde{r}^a
\]

(2.28)

as mandated by spherical symmetry. Note however that \( \vartheta \) does not have a simple analytical form,

\[
\vartheta = \frac{1}{8\pi\beta^2} \left[ Q(\beta\mu + \beta g A_0/2) - Q(\beta\mu - \beta g A_0/2) \right]
\]

\[
Q(u) = 2 \int_0^\infty ds \ln \left[ \frac{e^s + e^u}{e^s + e^{-u}} \right] \simeq \left( u^2 + \frac{\pi^2}{3} \right) \tanh \left( \frac{12 \ln 2}{\pi^2} u \right).
\]

(2.29)

where the analytic approximation to \( Q \) is accurate to about 0.62%; the derivative is accurate to 0.92%.

In the examples which we consider in detail we will be interested in the limit where \( A_0 \) is large and where \( \beta \to 0 \) or \( \mu \to 0 \). In these cases we have

\[
\vartheta \simeq \frac{g^2}{16\pi} A_0^2.
\]

(2.30)

We will require the volume charge density in the bulk to smoothly join the surface charge density at the surface layer, that is, \( \rho^a \lambda = \sigma^a \) at \( r = R \). Thus we
impose, \( q\lambda = \vartheta \) at \( r = R \) which, keeping in mind that the solutions produce large values of \( A_0 \) at \( R \), is equivalent to \( A_0^2/(4\pi) = \lambda A_0^3/(3\pi^2) \), or equivalently

\[
f(R) = \frac{3\pi R}{2\lambda}
\]

(2.31)

It is of course possible to modify this condition by requiring only that, at \( r = R \), \( q\lambda = s\vartheta \) for some number \( s = O(1) \), which is equivalent to replacing \( \lambda \rightarrow \lambda/s \); our results are insensitive to such a replacement.

To determine the consequences of the second of the above two conditions (color confinement) we need the components of the chromo-electric and chromo-magnetic fields parallel and perpendicular to \( r \),

\[
\hat{r} \cdot E^a = -\frac{1}{g^2\lambda} \left( \frac{f}{x} \right)' \hat{r}^a \quad \hat{r} \cdot B^a = -\frac{1}{g^2\lambda} \frac{a^2 - 1}{x^2} \hat{r}^a
\]

\[
(\hat{r} \times B^a)^i = -\frac{1}{g^2\lambda} \frac{a'}{x} \epsilon_{ija} \hat{r}^j \quad (\hat{r} \times E^a)^i = -\frac{1}{g^2\lambda} \frac{f a}{x^2} \epsilon_{ija} \hat{r}^j
\]

(2.32)

The first of these relations, together with the previously derived result \( (f/r)' \neq 0 \), implies that color will leak from the system unless an appropriate modification is included. The situation is identical to the one present in the bag model\(^9\), and the solution which we adopt is the same\(^10\). We will couple our system at the \( r = R \) boundary to a CP-odd field \( \eta' \) via a term proportional to the Chern-Simons term; this coupling insures that color is confined to the region \( r \leq R \). Denoting by \( F_{\eta'} \) the decay constant on the \( \eta' \), the coupling to this field at \( r = R \) are determined by the relations

\[
\hat{r} \cdot E^a = \frac{\alpha}{\pi F_{\eta'}} \hat{r} \cdot B^a \eta' \quad \hat{r} \times B^a = -\frac{\alpha}{\pi F_{\eta'}} \hat{r} \times E^a \eta'
\]

(2.33)

from which we derive \( (\hat{r} \cdot E^a)(\hat{r} \times E^a) + (\hat{r} \cdot B^a)(\hat{r} \times B^a) = 0 \); in terms of the \( a \)
and $f$ fields this becomes

$$fa(xf' - f) + xa'(a^2 - 1) = 0 \text{ at } r = R \quad (2.34)$$

which is the desired condition.#3

Character of the solutions

The TFQCD potentials $f$ and $a$ are then obtained by solving the equations (2.26) subject to the boundary conditions (2.27), (2.31) and (2.34). These solutions, as well as all thermodynamic variables, will depend on the parameter $\alpha = g^2/(4\pi)$. In order to specify $\alpha$ we first fix the thermodynamic variables of the system, such as the energy and volume; the TFQCD expresses these thermodynamic variables as functions of $\alpha$, which is chosen so that the chosen values are met.

When considering the (2.26) we find that, for given values of $X$ and $\alpha$, there are several solutions satisfying the boundary conditions#4. Of these solutions there is a set (we, in fact, found two such solutions) which minimizes $\Omega$, the thermodynamic potential at constant pressure and chemical potential

$$\Omega = - \int d^3x \, \mathcal{P}, \quad (2.35)$$

where $\mathcal{P}$ denotes the total pressure. Numerical studies show that there is no cross-over as $\alpha$ changes: each member of the set of solution which minimizes $\Omega$

---

#3 Concerning (2.34) we know, from the numerical integration of (2.26), that $f(xf' - f)$ does not vanish, we also find that it is numerically large for the situations we consider in detail. It follows that (2.34) can be approximately replaced by the simpler condition $a(R) = 0$.

#4 This is reminiscent of the situations found in the case of classical solutions to the Yang-Mills equations with external sources13).
is a smooth function of $\alpha$. Selecting the solution which minimizes $\Omega$ we then determined $\alpha$ by matching the specified energy and baryon number.

The explicit expressions for $\Omega$, the total energy $E$ and the total number of particles (baryon number) $N$ are

\[
\Omega = \frac{1}{\alpha \lambda} \int_0^X dx \left\{ \frac{1}{2} \left( \frac{f^2}{x} - \frac{f}{x} \right)^2 + \frac{1}{2} \left( \frac{1-a^2}{x} \right)^2 - \left( \frac{f a}{x} \right)^2 - \frac{\alpha}{24 \pi} f^2 (2 + \frac{f^2}{x^2}) \right\} - \frac{2}{9 \pi} \left[ \mu^4 + 2 \mu^2 (\pi k T)^2 + \frac{13}{15} (\pi k T)^4 \right] R^3
\]

\[
E = \frac{1}{\alpha \lambda} \int_0^X dx \left\{ \frac{1}{2} \left( \frac{f^2}{x} - \frac{f}{x} \right)^2 + \frac{1}{2} \left( \frac{1-a^2}{x} \right)^2 + \left( \frac{fa}{x} \right)^2 + \frac{\alpha}{24 \pi} f^2 (2 - \frac{f^2}{x^2}) \right\} + \frac{2}{3 \pi} \left[ \mu^4 + 2 \mu^2 (\pi k T)^2 + \frac{13}{15} (\pi k T)^4 \right] R^3
\]

\[
N = \frac{2 \mu}{\pi} \left\{ \frac{4}{9} \left[ \mu^2 + (\pi k T)^2 \right] R^3 + \lambda \int_0^X dx \ f^2 \right\},
\]

where

\[ X = \frac{R}{\lambda}. \]

For future reference we also provide the expression for the (total) entropy of the system

\[
\frac{1}{k} S = \frac{2}{3} \lambda \pi k T \left\{ \frac{4}{3} X^3 \lambda^2 \left[ \mu^2 + \frac{13}{15} (\pi k T)^2 \right] + \int_0^X f^2 \ dx \right\}
\]

Solutions of the equations for $f$ and $a$ can be obtained using standard numerical algorithms; due to the singular nature of the equations at the origin the relaxation method is best suited.
Fig. 1 Examples of solutions $f(x)$ (dashed lines) and $a(x)$ (solid lines) corresponding to $R = 10$ fm, $E/V = 4$ GeV/fm$^3$ and $T = 150$ MeV ($X = 47$). Cases (a) and (b) corresponds to the solutions which minimize the thermodynamic potential. Cases (c) and (d) have larger $\Omega$ and represent unstable solutions; the values of $\alpha$ corresponding to each solution are indicated.

We present several examples of the solutions in Fig. 1 where we took $R = 10$ fm, $E/V = 4$ GeV/fm$^3$ and $T = 150$ MeV (which implies $X = 47$). All the solutions in Fig. 1 satisfy the equations (2.26) and the boundary conditions; the solutions which minimize the thermodynamic potential corresponds to cases (a) and (b)\#5.

\#5 The singular nature of the equations allows for the multiplicity of solutions; we have found 8 solutions in total (for the given values of $E$ and $V$ but having different values of $\alpha$), though we cannot assert that this an exhaustive list. Using the relaxation method, the solution that minimizes $\Omega$ was least sensitive to the initial trial functions, solutions with larger $\Omega$ become increasingly more difficult to find as the range of initial configurations which relax to such solutions of (2.26) becomes more and more restricted. We have not attempted to perform an complete study of the properties and number of solutions restricting ourselves to finding the one solutions relevant for physical applications together with some examples of unstable solutions.
These two solutions correspond to indistinguishable thermodynamics (within numerical errors); for the calculations below this duality presents no complications. We have not attempted to study the stability of these solutions against non-radially symmetric perturbations \(^20\).

Given these results we must now determine whether they are consistent with the original assumptions, that is, whether \(f\) varies slowly enough to be considered constant in a region of width \(\sim \lambda\). We also must determine to what extent are color charges screened. The plots presented correspond to both cases (a) and (b) in Fig. 1.

![Fig. 2](image)

**Fig. 2** Validity of the Thomas-Fermi approximation. (a): only in the regions near the boundary at \(x = 47\) \((x > 45.5)\) and the origin \((x < 1)\) the approximation does breaks down. (b): charge is effectively screened throughout the volume (the logarithm is base 10).

The rate of change of \(f\) is sufficiently slow provided the potential \(A_0\) changes little within a region of size \(\lambda\), this is equivalent to

\[
\frac{(f/x)'}{(f/x)} < 1; \quad (2.41)
\]

a plot of the left hand side of this equation if presented in Fig. 2(a). We see that the condition (2.41) is satisfied except in the vicinity of the origin and the
$r = R$ boundary. The value of (2.41) near $x = 0$ presented in Fig. 2(a) is an underestimate generated by numerical errors (the equations are singular at $x = 0$); for $x \to 0$, $(f/x)^{\prime}/(f/x) \simeq 1/x$.

The magnitude of the charge in a subvolume $\delta V \sim \lambda^3$ is obtained from (2.14), and equals

$$
\delta q = \lambda^3 q = \frac{1}{24\pi^2} \frac{f}{x} \left( \frac{f^2}{x^2} + 1 \right),
$$

(2.42)

A plot of this quantity is presented in Fig. 2(b); as can be seen the magnitude of the color charge inside each subvolume is quite small except near the $r = R$ boundary: the system does screen its charges quite effectively.

2.5. Solutions for Small $X$

When $X$ is small then $f$ will be small also since it is monotonic (this follows from the boundary condition (2.31)). In this case the equation for $a$ decouples and so does the boundary condition (2.34),

$$
x^2 a'' + (1 - a^2)a = 0; \quad a(0) = 1, \ a'(X) \left[ a(X)^2 - 1 \right] = 0
$$

(2.43)

If we define

$$
a_2 = \frac{1}{2} a''(0),
$$

(2.44)

it is easy to see that the solution to the above equation is a function of $a_2 x^2$. It is then enough to assume $a_2 = \pm 1$; the general solutions are obtained from these by rescaling $x$. The solutions to the above differential equation (for $a_2 = \pm 1$) are
presented in Fig 3. The solutions are monotonic, so the boundary condition at \( x = X \) is satisfied when \( a(X) = -1 \) which occurs only for \( a_2 < 0 \), numerically

\[
a_2X^2 \simeq -4.1
\]

which completely specifies the solution.

\[\text{Fig. 3} \quad \text{Solutions for small } X, \text{ solid curve: } a''(0) = -1, \text{ dashed curve: } a''(0) = 1.\]

Again neglecting \( f \) and evaluating numerically the integrals gives

\[
\mathcal{E} \simeq \frac{5.42}{\alpha R} + 4.28(k\mathcal{T})^4\mathcal{V}
\]

\[
\mathcal{P} \simeq \frac{2.91}{\alpha} \mathcal{V}^{-4/3} + 1.43(k\mathcal{T})^4
\]

For example, at \( \mathcal{T} = 0, \mathcal{E} = 1 \text{ GeV}, R = 1\text{ fm}, \alpha \simeq 1 \) and \( \mathcal{P} \simeq 85 \text{ MeV/fm}^3 \). Fixing \( \mathcal{V} \) and \( \mathcal{T} \), the coupling strength \( \alpha \) drops as \( 1/\mathcal{E} \).

At zero temperature we have \( \mathcal{N} \simeq 2\mu^3\mathcal{V}/(3\pi^2) \) and

\[
\mathcal{E} \simeq \frac{5.42}{\alpha R} + \frac{\mu^4\mathcal{V}}{19.74}, \quad \mathcal{P} = \frac{2.91}{\alpha} \mathcal{V}^{-4/3} \frac{\mu^4}{59.04}
\]

so that the equation of state becomes

\[
\mathcal{P} \mathcal{V}^{4/3} \simeq \frac{2.91}{\alpha} + \frac{\mathcal{N}^{4/3}}{1.62}
\]
In this case $N = 2$, $E = 1$ GeV and $R = 1$ fm imply $\alpha \simeq 2.4$ and $P \simeq 77$ MeV/fm$^3$.

The numbers obtained for the case of small $X$ are then quite consistent with those obtained using the bag model\(^9\) (except perhaps for a large value for $\alpha$). Note however that in the present calculation the contributions from the non-ideal gas terms are very important and the numerical agreement is not trivial. It is also true that the present model is far from realistic (being based on an $SU(2)$ gauge theory with a single species of massless quarks). These results are therefore quite encouraging but not conclusive as to the physical relevance of this model.

Concerning the other thermodynamic quantities they relapse to their free-particle values up to $O(f^2)$ corrections. Note that the adiabats are, in general, defined by $X = \text{const.}$ which, for the case $N = 0$, imply $P^3 V^4 = \text{const.}$ just like a relativistic ideal gas.

When $\mu = 0$ an approximate solution for $f$ which satisfies the boundary conditions is

$$f \simeq 4.57 \frac{x^2}{X},$$

in this case the entropy becomes $S \simeq 1.5kX^3$ and the heat capacity equals $C_V \simeq 3S$; the largest contribution to these quantities ($\sim 94\%$) comes from the $\int f^2$ term.

2.6. Solutions for large $X$

In order to study the solutions to (2.26) for $x$ finite but $X \to \infty$ it proves convenient to define $y = x/X$. We are then interested in the small $y$ behavior of
the solutions and a power series is appropriate,

\[ f = f_2 \left[ y^2 + \left( \frac{2a_2}{5} + u \right) y^4 + \left( \frac{6a_2^2}{35} - \frac{f_2^2}{70} + \frac{5uf_2^2}{14X^2} + \frac{2ua_2}{7} + \frac{14u^2}{5} \right) y^6 + \cdots \right] \]

\[ a = 1 + a_2y^2 + \left( \frac{3a_2^2 - f_2^2}{10} \right) y^4 + \left( \frac{a_3}{10} - \frac{3a_2f_2^2}{35} - \frac{uf_2^2}{14} \right) y^6 + \cdots \]

(2.50)

where \( u = \alpha X^2/(60\pi) \). Numerical simulations indicate that neither \( f_2 \) nor \( a_2 \) increase with \( X \) which, using (2.31), (2.34), leads to \( f \sim 3\pi x^2/2X^2 \) and \( a \sim 1 - x^2/X^2 \) for \( x \ll X \). Thus, as \( X \to \infty \), \( f \to 0 \) and \( a \to 1 \) for \( x \) finite.

For \( x \ll X \) the boundary conditions require \( a \simeq 0 \) and \( f \gg 1 \); the first of the equations (2.26) can then be approximated by

\[ f'' \simeq \frac{\alpha}{6\pi} \frac{f^3}{X^2}; \quad x \ll X \]  

(2.51)

whose solution (using (2.31)) reads

\[ f \simeq \frac{6\pi X}{4 + (X - x)\sqrt{3\pi\alpha}}. \]  

(2.52)

Using these results we can evaluate the various thermodynamic quantities for large \( X \). For example,

\[ \mathcal{E} - \mathcal{E}_{\text{ideal gas}} \simeq \sqrt{\frac{3\pi}{16\alpha}} \frac{R^2}{\lambda^3}; \quad X \gg 1 \]  

(2.53)

from which we find \( |\mathcal{E} - \mathcal{E}_{\text{ideal gas}}|/\mathcal{E}_{\text{ideal gas}} \sim 1/X \). All other thermodynamic quantities exhibit this behavior: for large \( T \) and fixed \( R \) (corresponding to large \( X \)) the system approaches a mixture of ideal gases.
We emphasize that this is not a result of asymptotic freedom (when the running of the coupling is included the large $X$ behavior will acquire logarithmic corrections), but a property of the solutions to the differential equations. In the infinite volume limit the charges are screened which requires $A_0 = 0$ (see (2.42)).

It is also worth noticing that (2.53) explicitly displays the finite-volume corrections to the ideal gas results.

3. Applications

We now consider some applications of the above formalism. We first study a system with vanishing baryon number (corresponding to $\mu = 0$), and then consider the case of zero temperature.

3.1. Zero baryon number

This situation is believed to be of relevance in relativistic heavy ion collisions, such as those to be produced at RHIC$^{21}$, where, in the standard picture, the nuclei will go through one another leaving behind a region of hot quark-gluon plasma with zero baryon number$^{22}$.

The requirement $\mathcal{N} = 0$ in (2.38) corresponds to $\mu = 0$ which simplifies some of the expressions. In particular the only scales in the system are the temperature and the volume. The plot of the equation of state is given in Figs. 4 and 5.
Fig. 4 Equation of state within the Thomas-Fermi approximation for the $N = 0$ case. The graph displays the pressure as a function of the volume for several values of the temperature ($P$ in MeV$^4$, $V$ in fm$^3$, $T$ in MeV; the logarithms are base 10).

Fig. 5 Three-dimensional rendition of the equation of state for $N = 0$; ($P$ in MeV$^4$, $V$ in fm$^3$, $T$ in MeV; the logarithms are base 10).

We have determined $\alpha$ by requiring the solution to minimize the thermodynamic potential $\Omega$ when the energy density equals 4 GeV/$\text{fm}^3$ at $T = 150$ MeV, $R =$
10fm (which is consistent with the expectations for RHIC); in this case $\alpha \simeq 1.568$.

If we now allow the system to expand adiabatically, we can use the above expressions to obtain the relationship between $T$ and $R$ corresponding to this process. This isentropic transformation describes (in an admittedly oversimplified manner) the expansion of a quark-gluon plasma. The entropy is gotten from (2.40) by setting $\mu = 0$, the result is

$$
\frac{1}{k} S = \frac{13}{135} X^3 + \frac{1}{3} \int_0^X dx \ f^2(x). \tag{3.1}
$$

Since $S$ is a function of $X$ only (a consequence of having only two scales in the problem, $R$ and $T$), the equation for the adiabats is $X =$constant, or, equivalently $V T^3 =$const. corresponding to an adiabatic index $\gamma = 4$. Note that the $\int f^2$ term in $S$ modifies the usual free fermion gas relation $S \propto T^3$; the corrections are $\sim 20\%$ (which is smaller than the corresponding contributions in the case of small $X$, see section 2.5).

We can also easily determine the energy density for this isentropic process. From the expression for the total energy in (2.37) it follows that $E \lambda$ is a function of $X$ only (for the $\mu = 0$ case). It follows that at constant entropy $E$ scales as $T$. The energy density then will scale as $T / R^3 \propto T^4$, just as for an ideal gas of massless particles.

Using the expression for $S$ we obtain the heat capacity at constant volume,

$$
\frac{1}{k} C_V = X \left( \frac{\partial S}{\partial X} \right). \tag{3.2}
$$

In calculating this expression one must remember that the boundary conditions to the TFQCD equations depend on $X$, so that we should in fact write $f = f(X; x)$;
when the partial derivative is taken in (3.2), \( f \) must also be differentiated under the integral sign.

### 3.2. Zero Temperature

We now turn to the case of zero temperature; the dimensional quantities in the system are now \( \mu \) and \( R \). In this case all dimensionless quantities such as \( \mathcal{E}/\mu \) will be functions of \( R\mu \) only. The chemical potential is determined in terms of \( R \) and \( \mathcal{N} \) using (2.38) but this must be done numerically since the non-ideal gas term is significant and cannot be ignored. The plot of the equation of state for this case is given in Fig. 6.

![Equation of state plot](image.png)

*Fig. 6* The equation of state within the Thomas-Fermi approximation for the case zero temperature case; \( P \) in MeV\(^4\), \( V \) in fm\(^3\); the logarithms are base 10.

The equivalent contour plot for various values of \( N \) is presented in Fig. 7.
Fig. 7  Pressure as a function of volume at zero temperature, for various values of log $N$; ($P$ in MeV$^4$, $V$ in fm$^3$; the logarithms are base 10).

The equation of state (for the range of variables presented in figure 7) is well represented by the equation

$$P V^{4/3} = z(N); \quad \log z(N) \simeq 12.82 + 5.46 \left[ \log \left( \frac{N}{1.91 \times 10^4} \right) \right]^{1/4} \quad (3.3)$$

The $P V^{4/3}$ behavior is a result of simple scaling arguments and is therefore present here as well as for small $X$. In contrast, the $N$ dependence of the equation of state is radically different (cf. (2.48)).

We have also determined the chemical potential as a function of temperature and volume. The result is presented in Fig. 8.
As $T \to 0$ the entropy goes to zero linearly,

$$S \xrightarrow{T \to 0} \frac{k}{\sqrt{3} \mu} \left[ \frac{X_o}{3} \int_0^{X_o} dx f^2(x) \right]$$

since the fermionic contribution dominates in this limit; we then also have $C_V = S$. 

Fig. 8 The chemical potential as a function of volume and baryon number at zero temperature; ($P$ in MeV$^4$, $V$ in fm$^3$, $\mu$ in MeV; the logarithms are base 10).
4. Extensions of the method

The inclusion of more flavors is quite straightforward, the charges generated by each simply add. Possible computational difficulties arise when the fermion mass cannot be neglected (as is the case for the strange quark) for in this case a closed form for the fermionic partition function is not available. We will not pursue here this situation further as it involves no new concepts.

A more interesting extension is obtained by considering $SU(3)$ as the gauge group. In this case there are two important modifications. First, within each subvolume $\delta V$, though we still have $\bar{A}_0 =$constant and diagonal, this now implies $\bar{A}_0 = g(\mathcal{A}_0 \lambda_3 + \mathcal{B}_0 \lambda_8)/(2i)$. In general $\mathcal{B}_0 \neq 0$, so in this case we will have additional contributions depending on this new potential. The TFQCD equations are derived in the same way as for the $SU(2)$ case. Therefore the presence of the gauge field is summarized by the replacements

$$
\mu \rightarrow \mu + \frac{g}{2} \left( \mathcal{A}_0 + \mathcal{B}_0 / \sqrt{3} \right),
$$

$$
\mu + \frac{g}{2} \left( -\mathcal{A}_0 + \mathcal{B}_0 / \sqrt{3} \right),
$$

$$
\mu - g\mathcal{B}_0 / \sqrt{3},
$$

in $Z_0$ (Eq. (2.11)) The resulting fermionic partition function is then

$$
\ln Z_\psi = \frac{\beta \delta V}{4 \pi^2} \left[ \frac{g^4}{24} (\mathcal{A}_0^2 + \mathcal{B}_0^2)^2 + \frac{\mu g^3}{3 \sqrt{3}} \mathcal{B}_0 (3\mathcal{A}_0^2 - \mathcal{B}_0^2) 
+ \frac{g^2}{12 \lambda^2} (\mathcal{A}_0^2 + \mathcal{B}_0^2) + \mu^4 + 2\mu^2 (\pi k T)^2 + \frac{7}{15} (\pi k T)^4 \right].
$$

(4.2)
Using this result we obtain the charge densities

\[
q_3 = \frac{g A_0}{6\pi^2} \left[ \frac{g^2}{4} \left( A^2_0 + B^2_0 \right) + \sqrt{3} g\mu B_0 + \frac{1}{4\lambda^2} \right],
\]

\[
q_8 = \frac{g}{2\sqrt{3}\pi^2} \left[ \frac{g^2}{4} B_0 \left( A^2_0 + B^2_0 \right) + \frac{\sqrt{3}}{2} g\mu \left( A^2_0 - B^2_0 \right) + \frac{B_0}{4\lambda^2} \right].
\]

(4.3)

The second modification concerns the form of the spherically symmetric Ansatz for the background gauge potentials. For \( SU(3) \) a possible Ansatz takes the form (now including a contribution in the \( \lambda_8 \) direction)

\[
\bar{A}^0_a = A_0 \hat{r}_a; \quad (a = 1, 2, 3)
\]

\[
\bar{A}^i_a = \epsilon_{iaj} \hat{r}_j \left( \frac{1 + \varphi_2}{r} \right) + \left( \delta_{ia} - \hat{r}_i \hat{r}_a \right) \frac{\varphi_1}{r} + \hat{r}_i \hat{r}_a A_1; \quad (a = 1, 2, 3)
\]

(4.4)

\[
\bar{A}^0_8 = B_0.
\]

Note however that the choice of the \( SU(2) \) subgroup in which the potentials \( \bar{A}^\mu_a, \ (a = 1, 2, 3) \) reside is arbitrary, and that it costs no energy to change from one such subgroup to another; these degrees of freedom are included through a set of collective coordinates\(^2\). The full Ansatz we use is then (we define \( \bar{A}^\mu = \lambda_n \bar{A}^\mu_n/(2i) \) where the \( \lambda_n \) denote the usual Gell-Mann matrices)

\[
\bar{A}^\mu \to \bar{A}^\mu = U^\dagger \bar{A}^\mu U
\]

(4.5)

where the \( \bar{A}^\mu_n \) are given in (4.4) and \( U \) is a time dependent \( SU(3) \) matrix.

The Lagrangian for the background gauge fields then becomes

\[
\frac{1}{2} \text{tr} \bar{F}^2_{\mu\nu} \to \frac{1}{2} \text{tr} \bar{F}^2_{\mu\nu} + 2 \text{tr} F^0_{ij} \left[ \bar{A}_i, \mathbb{I} \right] + \text{tr} \left\{ \left[ \bar{A}_i, \mathbb{I} \right] \left[ \bar{A}_i, \mathbb{I} \right] \right\}
\]

(4.6)
where
\[ \mathbb{I} R = \dot{U} U^\dagger, \quad (4.7) \]
and \( \bar{F} \) is the field strength corresponding to \( \bar{A} \). When the form of the gauge potentials in the \( SU(2) \) subgroup takes the form \eqref{eq:4.4}, \( \mathbb{I} R \) should have no components along the generators of the \( SU(2) \) subgroup generated by \( \lambda_{1,2,3} \), that is we take
\[ \mathbb{I} R = \sum_{n=4}^{8} \frac{1}{2} \lambda_n \mathbb{I} R^n \quad (4.8) \]
which considerably simplifies \eqref{eq:4.6}. The corresponding action is
\[ S = \int d^4x \frac{1}{2g^2} \text{tr} F_{\mu \nu}^2 + \frac{1}{2} c^2 \int dt \text{tr} \dot{U}^\dagger \dot{U}; \quad c^2 = \frac{1}{\alpha} \int_0^R dr (a - 1)^2. \quad (4.9) \]
Numerically the coefficient \( c \) can be very large (for the numerical solutions presented \( c \sim 3.5 \times 10^3 \lambda \))

We will use the notation
\[ B_0 = \frac{h(r)}{rg}. \quad (4.10) \]
whence the TFQCD equations become (a prime denotes a derivative with respect to \( x = r/\lambda \)),
\[ f'' = \frac{2a^2}{x^2} f + \frac{\alpha}{6\pi} f \left[ \frac{f^2 + h^2}{x^2} + \left( 4\sqrt{3} \lambda \mu \right) \frac{h}{x} + 1 \right] \]
\[ a'' = \frac{a^2 - f^2 - 1}{x^2} \frac{a}{a} \]
\[ h'' = \frac{\alpha}{6\pi} \left[ \frac{h(f^2 + h^2)}{x^2} + \left( 2\sqrt{3} \lambda \mu \right) \frac{f^2 - h^2}{x} + h \right] \quad (4.11) \]
\[ \#6 \] The simplicity of this result is a consequence of the fact that \( U \) is made to reside in \( SU(3)/SU(2) \sim S^5 \), a five-dimensional sphere, where the number of invariants is very limited. The solutions of to the classical equations of motion for \( U \) are geodesics representing a motion along the great circles of \( S^5 \).
which can be solved using the same methods as before. Note that \( h = 0 \) is not allowed when \( \mu \neq 0 \).

For the interesting case \( \mu = 0, h = 0 \) is a solution to the above equations. Hence, for zero baryon number, the previous solutions also satisfy the \( SU(3) \) TFQCD equations. It does not follow, however, that these solutions again minimize the thermodynamic potential. Note also that even in the case \( h = 0 \) there is an additional contribution to the thermodynamic functions from the collective variables \( U \).

We will not pursue this case further in this paper. A realistic investigation of the \( SU(3) \) case requires we include mass term for the (strange) quarks, and also the contributions of the collective coordinates to the thermodynamics of the system. We will consider these issues in a forthcoming publication.

5. Conclusions

We have presented an approximate treatment of QCD based on the same ideas as the Thomas Fermi atom. Within this framework the thermodynamics of the system can be derived and the results can be compared with the experimental results which will soon be available.

The method is based on a subdivision of the system into subvolumes which are still large enough to be considered statistical systems. These subvolumes interact through an average gauge field whose sources are the thermodynamically averaged non-Abelian charges for the subvolumes. These charges, though small, are not completely screened due to the assumed smallness of the subvolumes.
The formalism was developed in this paper for the simplified case of an $SU(2)$ gauge group, though we did provide a brief discussion of the modifications required for and $SU(3)$ theory. We also ignored fermion masses and all interactions between the partons inside each of the subvolumes. Nonetheless we found that the numerical values for the pressure in the small $\mathcal{N}$ case are in rough agreement with the bag-model calculations.

For large temperatures, or densities ($X \gg 1$) the solutions to the equations of motion are such that all thermodynamic quantities approach those of a mixture of ideal gases, with $1/X$ measuring the deviation from this limiting behavior. This feature is not related to asymptotic freedom but a result of screening.

In the limit $R \to \infty$ we have $f = 0$ and $a = 1$, and the equation of state reduces to that of an ideal gas. This model then provides an approximation to the finite-volume corrections to the ideal gas, this is explicitly demonstrated in (2.53) which gives the surface corrections to the energy of the system.

A realistic calculations must be performed for an $SU(3)$ gauge theory with massive fermions; the partition function inside each subvolume should be evaluated to the highest order available (or possible) in perturbation theory. The inclusion of radiative corrections will induce, among other things, a dependence of the (now running) coupling constants on the temperature and chemical potential. For the present calculation no such effects were included. Finally one should also include finite volume effects as well as the corrections induced by the gluonic partition function. We will investigate such realistic situations in a forthcoming publication.

We found two solutions to the equations of motion satisfying the boundary
conditions and which minimize the thermodynamic potential $\Omega$. Both lead to the same thermodynamics and appear indistinguishable except near the origin (at least within numerical errors). A complete study of the behavior of these solutions under non-spherical perturbations along the lines of Ref 20 is required to determine the one which is most stable. We have not performed such an investigation since the presence of two such solutions does not alter the thermodynamics derived within the TFQCD approach.

The above treatment was not based on a semi-classical expansion of the partition function for the complete system. It is indeed possible to consider such an approach and use (2.12) as an approximation to the fermionic contribution. Then the integration over the gauge fields can be approximated by a saddle point method. We have not done this because the effective action which is to be minimized in the last step is, due to the Thomas-Fermi approximations used to obtain $Z_{\psi}$, unbounded from below. It is found that the solutions will soon violate the Thomas-Fermi conditions and the method is not consistent; this is displayed explicitly in the appendix for the case of QED. In contrast, the approach described in the above is consistent with the original approximations.

APPENDIX

In this appendix we present a semi-classical calculation of the partition function of QED using the Thomas-Fermi approximation for the fermionic partition function. The general expression is

$$Z = \int [dA][d\psi][d\bar{\psi}] e^{S} \quad (A.1)$$
where $S = S_g + S_\psi$, the first term denoting the gauge contribution, the second all terms involving the fermions. By definition we have

$$Z_\psi = \int [d\psi][d\bar{\psi}]e^{S_\psi} \quad (A.2)$$

which is approximated by $Z_\psi \simeq \int d^4x P_\psi$, where

$$P_\psi = \frac{1}{12\pi^2} \left[ (\mu + e\phi)^4 + 2(\pi kT)^2 (\mu + e\phi)^2 + \frac{7}{15} (\pi kT)^4 \right] \quad (A.3)$$

where $\phi$ denotes the electrostatic potential, and $e$ the charge of the fermions (only one flavor is considered). Note that $P_\psi$ is positive definite.

Assuming spherical symmetry the gauge potentials are of the form, $\phi = \phi(r)$, $A = a(r)r$. Choosing the $a = 0$ gauge gives the following expression for the partition function $Z = \int[d\phi]\exp S_{\text{eff}}$, where

$$S_{\text{eff}} = 4\pi \int_0^\beta dt \int_0^R dr \ r^2 \left[-\frac{1}{2} (\phi')^2 + P_\psi \right] \quad (A.4)$$

where $\beta$ denotes the inverse temperature, $R$ is the radius of the spherical vessel containing the system, a dash denotes a derivative with respect to $r$, the radial coordinate, and $t$ denotes the Euclidean time variable.

The integrand in $S_{\text{eff}}$ is not positive definite. Consider for example $\phi = \phi_0 \cos(kr + \nu)$ for constant $\nu$. If $\phi_0$ is sufficiently small and $k$ sufficiently large, the first term in $S_{\text{eff}}$ will dominate; the larger $k$, the more negative $S_{\text{eff}}$ becomes. The problem in this case is that these expressions for the scalar potential violate the Thomas-Fermi condition which requires $\lambda \phi' / \phi \ll 1$. This shows that a semi-classical treatment of the partition function is inconsistent with the Thomas-Fermi
approximation. We have verified that the same problems arise in the non-Abelian case.

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**FIGURE CAPTIONS**

1) Examples of solutions $f(x)$ (dashed lines) and $a(x)$ (solid lines) corresponding to $R = 10\text{fm}$, $E/V = 4 \text{GeV/fm}^3$ and $T = 150 \text{MeV}$ ($X = 47$). Cases (a) and (B) corresponds to the solutions which minimize the thermodynamic potential. Cases (c) and (d) have larger $\Omega$ and represent unstable solutions; the values of $\alpha$ corresponding to each solution are indicated.

2) Validity of the Thomas-Fermi approximation. Only in the regions near the boundary at $x = 47$ ($x > 45.5$) and the origin ($x < 1$) the approximation does breaks down.

3) Solutions for small $X$, solid curve: $a_2 = -1$, dashed curve: $a_2 = 1$. 

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4) Equation of state within the Thomas-Fermi approximation for the $\mathcal{N} = 0$ case. The graph displays the pressure as a function of the volume for several values of the temperature ($P$ in MeV$^4$, $V$ in fm$^3$, $T$ in MeV; the logarithms are base 10).

5) Three-dimensional rendition of the equation of state for $\mathcal{N} = 0$ ($P$ in MeV$^4$, $V$ in fm$^3$, $T$ in MeV; the logarithms are base 10).

6) The equation of state within the Thomas-Fermi approximation for the case zero temperature case; ($P$ in MeV$^4$, $V$ in fm$^3$; the logarithms are base 10).

7) The equation of state within the Thomas-Fermi approximation for the case zero temperature case; ($P$ in MeV$^4$, $V$ in fm$^3$; the logarithms are base 10).

8) The chemical potential as a function of volume and baryon number at zero temperature; ($P$ in MeV$^4$, $V$ in fm$^3$, $\mu$ in MeV; the logarithms are base 10).