Confinement may be more easily demonstrated at finite temperature using the Polyakov loop than at zero temperature using the Wilson loop. A natural mechanism for confinement can arise via the coupling of the adjoint Polyakov loop to $F_{\mu\nu}^2$. We demonstrate this mechanism with a one-loop calculation of the effective potential for $SU(2)$ gluons in a background field consisting of a non-zero color magnetic field and a non-trivial Polyakov loop. The color magnetic field drives the Polyakov loop to non-trivial behavior, and the Polyakov loop can remove the well-known tachyonic mode associated with the Saviddy vacuum. Minimizing the real part of the effective potential leads to confinement, as determined by the Polyakov loop. Unfortunately, we cannot arrange for simultaneous stability and confinement for this simple class of field configurations. We show for a large class of abelian background fields that at one loop tachyonic modes are necessary for confinement.

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I. INTRODUCTION

While the problem of quark confinement has been with us for some time [1], a completely satisfactory explanation of the mechanism has been elusive. It is widely held that unbroken non-Abelian gauge theories functions as dual superconductors, with some analog of the Meissner effect in type II superconductors confining chromoelectric flux into flux tubes. A fully satisfactory theory of quark confinement would be Lorentz invariant and gauge invariant, and have correct renormalization group behavior, which suggests that a successful theory must maintain a tight connection with perturbation theory.

Confinement is easiest to demonstrate at finite temperature, where it can be studied using the trace of the Polyakov loop \( \mathcal{P} = \mathcal{T} \exp \left[ i \oint_0^\beta dt A_0 \right] \). In pure gauge theories, without fundamental representation quarks, confinement is equivalent to the vanishing of the expected value of the trace of the Polyakov loop in the fundamental representation, \( Tr_F \mathcal{P} = 0 \). In contrast, demonstrating confinement at zero temperature requires showing that the Wilson loop has area law behavior for large areas. To use a simple analogy, it is much easier to study one-point functions than two-point functions. In pure \( SU(N) \) gauge theories, there is a global \( Z(N) \) symmetry, which, if unbroken, enforces the vanishing of the Polyakov loop. The spontaneous breaking of this \( Z(N) \) symmetry at high temperature is observed in lattice gauge theory simulations. In accord with these results, a successful theory of confinement should also show that the expected value of the trace of the Polyakov loop vanishes at low temperatures, and is non-zero at high temperatures.

The perturbative effective potential for the Polyakov loop, or equivalently, \( A_0 \) in an appropriate gauge does not indicate confinement [2,3]. The general result for \( SU(N) \) is a quartic polynomial in \( A_0 \). The minimum of this effective potential is always at \( A_0 = 0 \), corresponding to \( tr_F(\mathcal{P}) = N \); in the pure gauge theory without quarks, there are also equivalent solutions related by the global \( Z(N) \) symmetry. There is a simple physical reason behind this: when \( A_0 = 0 \), the free energy associated with the quark-gluon plasma is minimized.
One possible origin of confinement is that at low temperatures different regions of space are in different $Z_N$ phases, in such a way that $tr_F(P)$ averages to zero. This possibility is realized in the strong-coupling region of $Z(N)$ lattice gauge theories, where individual Polyakov loops can only take on values in $Z(N)$, but the average over an infinite volume is zero. This $Z(N)$ confining phase is connected to a phase of $SU(N)$ lattice gauge theories: if the Wilson action $\beta F S_F$ has added to it an adjoint term $\beta A S_A$, then the $Z(N)$ gauge theory can be obtained in the limit $\beta A \to \infty$ with $\beta F$ fixed. The region $\beta A$ large and $\beta F$ small realizes confinement in the same manner as the $Z(N)$ model. However, this region of the $\beta F-\beta A$ plane is separated from the continuum limit fixed point by a line of first-order phase transitions \[4\]. In contrast, examination of lattice field configurations generated by Monte Carlo simulation along the $\beta A = 0$ axis indicate that the magnitude of individual Polyakov loops is small at low temperatures and increases dramatically at the confinement transition. This argues against the possibility that confinement occurs because different regions of space are in different $Z_N$ phases.

We advocate another possibility: that the coupling between the local gauge field $F_{\mu\nu}$ and the adjoint Polyakov loop produces an effective action which leads to two different phases. In the low temperature phase, there is a magnetic condensate and the Polyakov loop indicates confinement; in the high temperature phase the magnetic condensate vanishes, the Polyakov loop indicates deconfinement, and the contribution of the thermal gauge boson gas dominates the free energy. This occurs because finite temperature effects from gluons naturally couple the local gauge field to adjoint representation Polyakov loops \[5\ 8\].

In order to produce confinement, the low temperature phase must minimize its free energy by minimizing the expected value of the trace of the adjoint Polyakov loop. For $SU(N)$, the fundamental and adjoint representation Polyakov loops are related by

$$Tr_A(P) = |Tr_F(P)|^2 - 1.$$ \hspace{1cm} (1.1)

Clearly, when $Tr_A(P)$ assumes its minimum value of $-1$, $Tr_F(P)$ assumes the value $0$. As seen in previous calculations, at sufficiently high temperature, contributions from thermal
gluons will dominate the free energy, driving the Polyakov loop towards its maximum value. In our scenario, this also drives the magnetic condensate to zero. The coupling of the gauge field to the adjoint Polyakov loop naturally incorporates the behavior known from lattice simulations: both the Polyakov loop and plaquette expectation values change at the deconfining phase transition of $SU(N)$ lattice gauge theories.

The idea that more than one field must be considered to describe the behavior of finite temperature QCD has previously been proposed by Campbell, Ellis and Olive [15,16]. In their effective theory, a phenomenological coupling of the gluon condensate to the chiral condensate was postulated. However, this approach does not include the important role played by the Polyakov loop in finite temperature chiral behavior. It is easy to show that there is a coupling between the chiral condensate and Polyakov loops. For example in references [5] and [6], we showed using lattice perturbation theory the role that the Polyakov loop plays in finite temperature corrections to the renormalization of the gauge field coupling constant. Differentiation of our result with respect to the quark mass shows perturbatively that the finite temperature corrections to the chiral condensate depend on the Polyakov loop as well as the gluon condensate. We have also shown in detail the interplay between the Polyakov loop and the chiral condensate for a variant of the Nambu-Jona-Lasinio model. [7,8]. A truly satisfactory treatment of the critical behavior of finite temperature QCD, including the effects of quarks, is likely to require treatment of the gluon condensate, the chiral condensate and the Polyakov loop.

II. CALCULATION OF THE FREE ENERGY

The observation of Saviddy [9] that the perturbative QCD vacuum is unstable with respect to the formation of a constant chromomagnetic field is one proposed origin for a magnetic condensate in QCD. However, Nielsen and Olesen [10] showed that such field configurations are themselves unstable, due to a tachyonic mode in the one-loop determinant. Several authors [11–14] have extended the work of Nielsen and Olesen to finite temperature,
treating the tachonic modes in various ways, including simply ignoring them.

We demonstrate the main points of our proposed mechanism by incorporating the effects of a non-trivial Polyakov loop, treating all modes exactly at one loop. The tachyonic mode can be removed by a non-trivial Polyakov loop. The free energy of an SU(2) gluon gas moving in a background field consisting of a constant color magnetic field and and a constant non-trivial Polyakov loop is obtained by calculating the one-loop, finite temperature effective potential in the background field gauge; see [17] and [18] for pedagogical introductions to the background field method. This is equivalent to summing up three contributions to the free energy: the classical action, the zero-point energies, and the free energy of the gluon gas. All finite temperature effects reside in the last term. As we have shown elsewhere, Polyakov loops appear naturally at one loop in an image expansion of finite temperature determinants [3–8]. Here, similar results are achieved by directly computing the functional determinant.

We choose the color magnetic field and the Polyakov loop to be simultaneously diagonal, and we let the color magnetic field \( H \) point in the \( x_3 \) direction. We take the Polyakov loop to be specified by a constant \( A_0 \) field, given in the adjoint representation by

\[
A_0 = \frac{\phi}{2\beta} \tau_3. \tag{2.1}
\]

The trace of the Polyakov loop is then given by

\[
Tr_F(\mathcal{P}) = 2 \cos(\phi/2) \tag{2.2}
\]

in the fundamental representation and by

\[
Tr_A(\mathcal{P}) = 1 + 2 \cos(\phi) \tag{2.3}
\]

in the adjoint representation. The external magnetic field we take to have the form

\[
A_2 = \frac{1}{2} H x_1 \tau_3 \tag{2.4}
\]

which gives rise to a chromomagnetic field.
As explained, for example, in [10], the external field gives rise to Landau levels in the gluon functional determinant.

The one-loop contribution to the free energy has the form [10,11]:

\[
2 \times \frac{1}{2} \sum_{m=0}^{\infty} \sum_{n, \pm} \frac{1}{\beta} \frac{H}{2\pi} \int \frac{dk_3}{2\pi} \ln \left[ \left( \omega_n - \frac{\phi}{\beta} \right)^2 + 2H \left( m + \frac{1}{2} \pm 1 \right) + k^2_3 \right] \\
+ \sum_n \frac{1}{\beta} \int \frac{d^3 \vec{k}}{(2\pi)^3} \ln \left( \omega^2_n + \vec{k}^2 \right),
\]

(2.6)

where the \( \omega_n = 2\pi n/\beta \) are the usual Matsubara frequencies, and where the terms \( 2H(m + 1/2 \pm 1) \) are the allowed Landau levels of the gauge field with its spin coupled to H. The explicit factor of 2 in front of the first term in the sum results from the trace over color degrees of freedom.

When \( \phi = 0 \), the \( n = 0 \) and \( m = 0 \) modes give rise to tachyonic modes for \( k_3 \) sufficiently small; these in turn give rise to an imaginary part in the free energy. We observe that these same modes will give a strictly real factor to the determinant provided

\[
\phi \geq \beta \sqrt{H}.
\]

(2.7)

However, if \( \phi \) is too large, it may cause the \( n = 1 \) mode to become unstable. Hence, the general criterion for stability is

\[
\beta \sqrt{H} < \phi < 2\pi - \beta \sqrt{H}
\]

(2.8)

A standard product representation [20]

\[
\prod_k \left[ 1 + \left( \frac{x}{2k\pi - a} \right)^2 \right] = \frac{\cosh(x) - \cos(a)}{1 - \cos(a)}
\]

(2.9)

allows us to write the effective potential in the form:

\[
V = \frac{1}{2g^2} H^2 + \frac{1}{2} \sum_{m=0}^{\infty} \sum_c \frac{H}{2\pi \beta} \int \frac{dk}{2\pi} \ln \{ \cosh [\beta \omega_{\pm}(m, k)] - \cos(\beta \phi) \} \\
+ \frac{1}{\beta} \int \frac{d^3 \vec{k}}{(2\pi)^3} \ln \left\{ \cosh \left[ \beta |\vec{k}| \right] - 1 \right\},
\]

(2.10)
where
\[ \omega_\pm^2(m, k) = 2H \left( m + \frac{1}{2} \pm 1 \right) + k^2 \] (2.11)

which can be negative. Note that the classical free energy of the magnetic field, \( H^2/2g^2 \), is included in this expression. Expanding the logarithm and discarding an irrelevant constant, the contribution of the Landau levels to the the effective potential can also be written in the form:
\[
\frac{1}{2} \sum_{m=0}^{\infty} \sum_{c} \frac{H}{2\pi} \int \frac{dk}{2\pi} \left[ \omega_\pm(m, k) - \frac{2}{\beta} \sum_{n=1}^{\infty} \frac{\cos(n\phi^c)}{n} e^{-n_\beta\omega_\pm(m, k)} \right].
\] (2.12)

After performing the sum over \( m \) in the equation above, we regulate the divergent portions of the resulting integral using the identity [10]
\[
(\omega^2 - i\delta)^{\nu-\mu} = \frac{i^{\mu-\nu}}{\Gamma(\mu-\nu)} \int d\tau \tau^{\mu-\nu-1} e^{-i\tau(\omega^2-i\delta)}. \tag{2.13}
\]

This yields a renormalized effective potential with real component
\[
V_R = \frac{11H^2}{48\pi^2} \ln \left( \frac{H}{\mu_0^2} \right) - \frac{(H)^{3/2}}{\pi^2 \beta} \sum_{n=1}^{\infty} \frac{\cos(n\phi)}{n} \left[ K_1(n\beta\sqrt{H}) - \frac{\pi}{2} Y_1(n\beta\sqrt{H}) \right] - \frac{2(H)^{3/2}}{\pi^2 \beta} \sum_{n=1}^{\infty} \frac{\cos(n\phi)}{n} \sum_{m=0}^{\infty} \sqrt{2m+3} K_1(n\beta\sqrt{2m+3}H) - \frac{2\pi^2}{90\beta^4} \tag{2.14}
\]

and imaginary component
\[
V_I = -\frac{H^2}{8\pi} - \frac{H^{3/2}}{2\pi \beta} \sum_{n=1}^{\infty} \frac{\cos(n\phi)}{n} J_1(n\beta\sqrt{H}), \tag{2.15}
\]

where \( \phi \) is defined in Eq. (2.1).

It is far from obvious that the potential we have found is real when the constraint of Eq. (2.8) is satisfied. We have checked numerically that the imaginary part of the potential behaves as it should. In figure 1, we show the imaginary part for \( \sqrt{H}/\mu = 0.5 \) and \( T/\mu = 0.25 \). The theory is stable in the region given by Eq. (2.8).

A similar calculation has been performed for different reasons by Starinets, Vshivtsev, and Zhukovskii [21]. Our results are in disagreement, with Bessel functions interchanged between the real and imaginary part. However, our results are in numerical agreement with
the exact result for $V_I$ derived by Cabo, Kalashnikov, and Shabad [22] for the case $\phi = 0$, and our result for $V_I$ is zero for the range of $\phi$ determined by Eq. (2.8). We are confident our results are correct.

Note that we have carefully chosen the overall additive constant of the free energy such that $V = 0$ when $H = 0$ and $T = 0$.

**III. MINIMIZATION OF RE(V)**

We are interested in showing that the system can lower its free energy by having $\phi = \pi$, or equivalently $Tr_F P = 0$. The simplest scheme is to minimize $Re(V)$ over all allowed values of $\phi$ and $H$, ignoring the imaginary part of $V$. If the global minimum of $Re(V)$ gives rise to an imaginary part, that will indicate that this field configuration is unstable, however.

Some insight into this issue can be obtained by noting that for very low temperatures, the finite temperature contributions to $Re(V)$ will be dominated by the $Y_1(n\beta\sqrt{H})$ terms. Using the asymptotic form for large argument, we see that they will favor the adjoint Polyakov loop being at its minimum, and hence favor confinement.

We have examined this issue numerically and found that at low temperatures, minimization of $Re(V)$ leads to $\phi = \pi$ being preferred. This is shown in figure 2, which plots the real part of the effective potential versus $H/\mu^2$ at $T/\mu = 0.25$ for both $\phi = 0$ and $\phi = \pi$. Figure 3 plots $V_I$, the imaginary part of the effective potential for the same parameters. Unfortunately, examination of $V_I$ shows that the lowest minima is not stable, lying just to the right of the stable region.

The Savvidy vacuum is preferred over the perturbative vacuum only for sufficiently low temperatures. Figure 4 shows the behavior of $V$ at $T/\mu = 0.73$ and $\phi = \pi$ compared with $V_{0T}$, the one-loop effective potential for a gluon gas at $H = 0$ and $A_0 = 0$. Above the pseudocritical temperature at approximately $T_c/\mu = 0.73$, the perturbative solution with $H = 0$ and $\phi = 0$ has lower free energy. Naively, this would indicate a first-order phase transition. However, the magnitude of $H/\mu^2$ is so large that a one-loop calculation cannot be
considered reliable. Further, examination of $V_I$ shows that the constant field configuration which minimizes $V$ is always unstable. All that can be said is that at low temperatures, the perturbative solution is definitely not the state of lowest free energy.

It is amusing to note that the behavior at very low temperatures reveals the existence of a familiar phenomenon associated with Landau levels, the de Haas-van Alphen effect. Fig. 5 plots the free energy as a function of $H/\mu^2$ at $T/\mu = 0.1$ for both $\phi = 0$ and $\phi = \pi$. The minima are the signs of the characteristic oscillatory components of the magnetic susceptibility.

Because of the possibility of fixing the external field $H$ in lattice simulations, it may be of interest to examine the analytical behavior at fixed $H$. At fixed $H$, $\phi = \pi$ is not necessarily the minimum value of $V$. This is demonstrated clearly in Fig. 6 and Fig. 7, which show the $V$ and $V_I$ as a function of $\phi$ at $T/\mu = 0.25$ for $H/\mu = 0.40, 0.65$ and $0.90$. For $H/\mu^2 = 0.40$, there is a region of $\phi$ where $V_I$ vanishes. The minimum of $V$ lies at or very near the boundary where $V_I$ develops an imaginary part. This particular case appears similar to the behavior found by van Baal, who examined self-dual solutions on a hypercube in the presence of a non-trivial Polyakov loop [25].

**IV. GENERALIZATIONS**

As we have seen above, the terms in the free energy which favor confinement come from the contribution of the would-be tachyonic modes to the free energy. We can generalize this result slightly. Consider an arbitrary time-independent vector potential of the form

$$\vec{A}(\vec{x})\tau_3$$

(4.1)

We write the eigenvalues that enter into the functional determinant in the form

$$\left[\left(\omega_n - \frac{\phi}{\beta}\right)^2 + \omega^2\right]$$

(4.2)

where $\omega^2$ can be positive or negative.
We write the positive eigenvalues as $\epsilon^2$, and the negative eigenvalues as $-\chi^2$. Following the identical reasoning used above, the positive eigenvalues contribute to the free energy a term

$$V_+ = \sum_n \frac{1}{\beta} \int d\epsilon \rho_+(\epsilon) \ln \left[ \left( \omega_n - \frac{\phi}{\beta} \right)^2 + \epsilon^2 \right],$$

where $\rho_+(\epsilon)$ is the eigenvalue density. It is easy to see that such eigenvalues always favor $\phi = 0$. On the other hand, the negative eigenvalues contribute to the free energy a term

$$V_- = \sum_n \frac{1}{\beta} \int d\chi \rho_-(\chi) \ln \left[ \left( \omega_n - \frac{\phi}{\beta} \right)^2 - \chi^2 \right] = C + \int d\chi \rho_-(\chi) \ln \left[ \cos^2 (\beta \chi) - \cos^2 (\phi) \right],$$

and such terms may favor $\phi \neq 0$. Thus in this simple generalization we have an indication of the likely importance of the negative modes in bringing about confinement.

It is interesting to note that the hard thermal loop resummation of Braaten and Pisarski drops precisely those features which we are advocating as the origin of confinement at low temperatures. The tachyonic modes in the Saviddy vacuum arise at the one-loop level from the spin coupling of the gauge field to the external magnetic field. However, the term accounting for this spin coupling is discarded in the hard thermal loop approximation, since the term is manifestly soft.

**V. CONCLUSIONS**

We have seen that finite temperature gauge theory is the best place to prove confinement, because one need only show that the expected value of the Polyakov loop vanishes. Models and mechanisms for which this cannot be demonstrated are not good candidate explanations of confinement. We have demonstrated these considerations with a one-loop calculation of the effective potential for $SU(2)$ gluons in a special background field consisting of a non-zero color magnetic field and a non-trivial Polyakov loop. The color magnetic field drives the Polyakov loop to non-trivial behavior, and the Polyakov loop can remove the well-known tachyonic mode associated with the Saviddy vacuum. Minimizing the real part of the
effective potential leads to confinement, as determined by the Polyakov loop. Unfortunately, we cannot arrange for simultaneous stability and confinement in this simple model. We are able to show for a class of abelian field configurations that tachyonic modes are necessary for confinement.

The mechanism we have described here for pure gauge theories is also applicable when the effects of quarks are included. The quark determinant has several effects, the most important of which is the explicit breaking of $Z(N)$ symmetry. Pure $SU(N)$ gauge theories have an exact global $Z(N)$ symmetry which, if unbroken, requires that the expectation value of $Tr_F P$ vanish. Fundamental representation particles, such as quarks, explicitly break this symmetry. However, explicit perturbative calculations shows that the symmetry breaking terms are numerically small, even when the current quark mass is zero. Strong-coupling arguments suggest that in the low temperature phase, it is the constituent quark mass rather than the current mass which controls the strength of the $Z(N)$ symmetry breaking, which would further reduce the effect.

Other effects of quarks enter from the change in renormalization group beta function coefficients, and from the increase in the number of degrees of freedom in the high temperature phase. At low temperatures, we expect that the competition between the gluonic terms which drive the trace of the Polyakov loop towards zero and the quark terms which drive it away from zero will produce a low temperature region where the Polyakov loop is small, in accord with lattice simulations.

It is obvious that the ground state of QCD is not well modeled by a constant chromomagnetic field. However, we can extend the insight we have gained into the possible origins of confinement by examining the confining properties of other background field configurations. There are several interesting classes of configurations, including constant field potentials. If the different space time components do not commute, this leads to constant color fields which are inequivalent to the class studied here. Another interesting class of background field configurations is abelian monopole configurations. It would also be very interesting to know if a suitably chosen ensemble of random background field configurations confines.
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FIGURES

FIG. 1. \( V_I \) as a function of \( \phi \) for \( H^{1/2}/\mu = 0.5 \) and \( T/\mu = 0.25 \).

FIG. 2. \( V \) at \( T/\mu = 0.25 \) for \( \phi = 0 \) and \( \phi = \pi \).

FIG. 3. \( V_I \) at \( T/\mu = 0.25 \) for \( \phi = 0 \) and \( \phi = \pi \).

FIG. 4. \( V \) at \( T/\mu = 0.1 \) for \( \phi = 0 \) and \( \phi = \pi \).

FIG. 5. \( V \) at \( T/\mu = 0.73 \) for \( \phi = \pi \) compared with \( V_{0T} \).

FIG. 6. \( V \) as a function of \( \phi \) at \( T/\mu = 0.25 \).

FIG. 7. \( V_I \) as a function of \( \phi \) at \( T/\mu = 0.25 \).
Figure 1: $V_I$ as a function of $\phi$ for $H^{1/2}/\mu = 0.5$ and $T/\mu = 0.25$
Figure 2: $V$ at $T/\mu = 0.25$ for $\phi = 0, \pi$
Figure 3: $V_I$ at $T/\mu = 0.25$ for $\phi = 0, \pi$
Figure 4: $V$ at $T/\mu = 0.1$ for $\phi = 0, \pi$
Figure 5: $V$ at $T_c/\mu = 0.73$ with $\phi = \pi$ compared to $V_{0T}$
Figure 6: $V$ as a function of $\phi$ at $T/\mu = 0.25$
Figure 7: $V_I$ as a function of $\phi$ at $T/\mu = 0.25$