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TEMPERATURE DEPENDENCE OF CRITICAL MAGNETIC FIELDS FOR THE ABELIAN HIGGS MODEL*

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ABSTRACT

One loop temperature and external electromagnetic field effects on the Abelian Higgs model are studied using the momentum space heat kernel. We obtain expressions for the critical fields necessary for symmetry restoration at some finite temperature and display the critical B vs. T curve separating the broken and restored phases in the B-T plane.

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I. INTRODUCTION

It is well known from the phenomenological theory of superconductivity [1] that the ordered or superconducting phase is destroyed by raising the temperature of the specimen above some critical value $T_c(0)$* or by applying an external magnetic field exceeding the critical value $B_c(0)$. In general, the critical B vs. T line in the B-T plane is described by the curve [2] in Fig. 1. From this, it is clear that at some finite temperature T (less than $T_c(0)$) the value of the critical field $B_c(T)$ is less than the zero temperature critical value $B_c(0)$, though the particular elliptic shape of the curve implies that the reduction is appreciable only for temperatures T close to $T_c(0)$. The interesting question which the present note examines, concerns the determination of the shape of the analogous curve in the case of the Abelian Higgs model which is just the relativistic analogue of the phenomenological Ginzburg-Landau theory of superconductivity. The hope being that perhaps a combination of finite temperature and external magnetic field effects will allow restoration of the broken symmetry (i.e., the transition to the normal or disordered state) to take place at lower values of B than those needed at zero temperature. The resulting B-T curve is displayed in Fig. 2.

II. ONE-LOOP EFFECTS

In this paper we shall be concerned specifically with the evaluation of one loop, finite temperature effects [3] for the massless abelian Higgs model coupled to an external electromagnetic field [4,5]. This system is characterized by the lagrangian (we choose the signature (-+++) for the metric):

$$\mathcal{L} = -\frac{1}{4} F_{\mu\nu}F^{\mu\nu} - \frac{1}{2}(D_\mu\varphi)_a(D^\mu\varphi)_a - \frac{\lambda}{4!}\varphi^4 \quad ; \quad a = 1,2 . \tag{1}$$

where

$$(D_\mu)_{ab} = \partial_\mu \delta_{ab} + e A_\mu \epsilon_{ab} \tag{2}$$

* In general $T_c(B)$ denotes the value of the critical temperature for some applied external field. Similarly we shall speak of the critical magnetic field $B_c(T)$.

and where the two real scalar fields φ_1 and φ_2 can also be conveniently written as one complex scalar field; $\varphi = \varphi_1 + i\varphi_2$ to display the similarity between φ and the order parameter of the Ginzburg-Landau theory. We shall use dimensional regularisation and renormalisation by minimal subtraction [6] when necessary.

The one loop effective action for this theory coupled to an external electromagnetic field $(F^{\mu\nu}_{cl.})$ is given in n-dimensions by:

$$\Gamma[F^R_{\mu\nu}, \varphi^R] = \int d^n x \left(\mathcal{L} + \frac{1}{2} F^R_{\mu\nu} F^{\mu\nu}_{cl.} \right) - \frac{1}{2} \ln \det \mathfrak{D} + \ln \det \mathfrak{D}_G \qquad (3)$$

where the last two terms are evaluated in Euclidean space with signature (++++). Here, the superscript R denotes renormalised quantities. The coupling to the external electro-magnetic field is given by the second term on the right hand side, while the objects $\mathfrak{D}$ and $\mathfrak{D}_G$ are matrices of various second functional derivatives of the action.

When one is interested in simply calculating one loop effects for the zero temperature theory, the configuration space heat kernel method [7] can be used with great advantage to calculate the functional determinants appearing in (3). This has been illustrated beautifully in the recent work of Shore [5]. We outline a generalisation of the technique enabling finite temperature calculations to be carried out, with the aid of a simple example.

Consider the case when $\mathfrak{D} = (-\partial^2 + m^2)\,\delta(x, y)$. The heat kernel for this propagator is easy to evaluate and is given by

$$\mathcal{G}(x,y;s) = \frac{s^{-\frac{n}{2}}}{(4\pi)^{\frac{n}{2}}}\, e^{-\frac{(x-y)^2}{4s}}\, e^{-m^2 s} . \qquad (4)$$

One simple technique of solving for this is to use the fact that

$$\mathcal{G}(x,y;0) = \delta(x,y) = \lim_{s\to 0} \frac{s^{-\frac{n}{2}}}{(4\pi)^{\frac{n}{2}}}\, e^{-\frac{(x-y)^2}{4s}} \qquad (5)$$

and then making the ansatz that

$$\mathcal{G}(x,y;s) = \frac{s^{-\frac{n}{2}}}{(4\pi)^{\frac{n}{2}}}\, e^{-\frac{(x-y)^2}{4s}}\, f(s) \qquad (6)$$

with $\lim_{s\to 0} f(s) = 1$.

The results of this simple example will be needed when we come to the evaluation of functional determinants for the particular $\mathfrak{D}$'s we are interested in, the only difference being that in general, m will depend on the constant field φ.

Returning to the example, and writing the Fourier transform of $\mathcal{G}(x, y;s)$ as $\mathcal{G}(k;s)$, we have

$$\begin{aligned} \ln \det \mathfrak{D} &= -\int_0^\infty \frac{ds}{s} \operatorname{Tr}_{x,y} \mathcal{G}(x,y;s) \\ &= -\int_0^\infty \frac{ds}{s} \int \frac{d^n k}{(2\pi)^n}\, \mathcal{G}(k;s) \int d^n x \, . \end{aligned} \qquad (7)$$

Finite temperature effects are now incorporated through the following prescription:

$$\int \frac{d^n k}{(2\pi)^n} \longrightarrow \frac{1}{\beta} \sum_{m=-\infty}^{+\infty} \int \frac{d^{n-1} k}{(2\pi)^{n-1}} \quad , \quad k_0 \longrightarrow \frac{2\pi m}{\beta} . \qquad (8)$$

Notice that the prescription differs from the usual one in the absence of factors of i. This is due to the fact that the phase space we work with is also Euclidean. After having introduced temperature, we should perhaps rename the non-thermodynamic quantities by their appropriate thermodynamic analogues. For example, the effective potential now becomes the free energy.

For the heat kernel given in equation (4), we write

$$\begin{aligned} \mathcal{G}(k;s) &= \int d^n x \; e^{-ik\cdot(x-y)}\, \mathcal{G}(x,y;s) \\ &= e^{-(k^2+m^2)s} , \end{aligned} \qquad (9)$$

so that

$$\ln \det \mathfrak{D} = -\int_0^\infty \frac{ds}{s} \left\{ \frac{1}{\beta} \sum_{p=-\infty}^{+\infty} e^{-\left(m^2 + \frac{4\pi^2 p^2}{\beta^2}\right)s} \; (4\pi s)^{-\frac{(n-1)}{2}} \right. \qquad (10)$$

after using prescription (8) and carrying out the $d^{n-1}k$ integration.

As it stands, it is not immediately clear how to extract the correct zero temperature result from this expression. However, the particular form of the summation we have to do can be written in terms of theta functions and the following, very useful identity derived [8];

$$\sum_{p=-\infty}^{+\infty} e^{-\frac{\pi^2 p^2}{s}} = \sqrt{\frac{s}{\pi}} \sum_{p=-\infty}^{+\infty} e^{-p^2 s} \qquad (11)$$

Using this, we can now write

$$\ln \det \mathcal{D} = -\sum_{p=-\infty}^{+\infty} \int_0^\infty \frac{ds}{(4\pi)^{\frac{n}{2}}} s^{-1-\frac{n}{2}} e^{-m^2 s} e^{-\frac{p^2\beta^2}{4s}} \qquad (12)$$

Now it is straightforward to separate this into a zero temperature piece (coming from the $p = 0$ term in the summation) and a temperature dependent piece. Thus,

$$\begin{aligned} \ln \det \mathcal{D} = & -\int_0^\infty \frac{ds}{(4\pi)^{n/2}} s^{-1-\frac{n}{2}} e^{-m^2 s} \\ & -2\sum_{p=1}^{\infty} \int_0^\infty \frac{ds}{(4\pi)^{n/2}} s^{-1-\frac{n}{2}} e^{-m^2 s} e^{-\frac{p^2\beta^2}{4s}} \end{aligned} \qquad (13)$$

Here, the first term is precisely the zero temperature contribution and needs to be regularised, while the second term is finite, as it should be, and we encounter no difficulties in letting $n \to 4$. The resulting integral is evaluated using the following formula [9]:

$$\int_0^\infty x^{\nu-1} e^{-\frac{\alpha}{x} - \gamma x} dx = 2\left(\frac{\alpha}{\gamma}\right)^{\nu/2} K_\nu\left(2\sqrt{\alpha\gamma}\right), \quad \mathrm{Re}\,\alpha, \gamma > 0 . \qquad (14)$$

The final expression, for large temperatures can be written down as

$$\ln \det \mathcal{D} = (T=0 \text{ result}) \int d^n x + \left(\frac{m^2}{12\beta^2} - \frac{\pi^2}{45\beta^4}\right) \int d^4 x + O(\beta^2), \qquad (15)$$

having used the asymptotic expansion for $K_2(z)$, the modified Bessel function.

III. RESULTS

We are now ready to compute the external electromagnetic effects at finite temperature using the method outlined above. We shall compute the energy densities in the broken and the restored phases. For the broken phase, the computation of finite temperature effects is a simple matter using the result

given in equation (15). For the restored phase however, we need to calculate the finite temperature corrections coming from the heat kernel of the operator $-D^2$ where $D_\mu = \partial_\mu - ieA_\mu$. It is given explicitly, for the case when $F_{\mu\nu}$ satisfies the constraint $\partial_\mu F_{\mu\nu} = 0$, by [5]

$$G(x,y;s) = \frac{s^{-\frac{n}{2}}}{(4\pi)^{n/2}} \Phi(x,y) \exp\left\{-\frac{1}{4}(x-y)_\mu (E\cot E)_{\mu\nu} (x-y)_\nu - \frac{1}{2}\mathrm{tr}\ln E^{-1}\sin E\right\} \qquad (16)$$

where E denotes the matrix (eFs) and $\Phi(x, y)$ is the path dependent phase factor

$$\Phi(x,y) = \exp i\int_y^x A_\mu(z)\, dz_\mu \qquad (17)$$

integrated along a straight line path from y to x.

A somewhat lengthy calculation [10] yields the following simple result for the case of a pure magnetic field directed along the z-axis:

$$\begin{aligned} -\frac{1}{2}\ln\det(-D^2) = & \int_0^\infty \frac{ds}{(4\pi)^{n/2}} s^{-1-\frac{n}{2}} e^{-\frac{1}{2}\mathrm{tr}\ln(E^{-1}\sin E)} \\ & + 2\sum_{p=1}^{\infty} \int_0^\infty \frac{ds}{(4\pi)^{n/2}} s^{-1-\frac{n}{2}} e^{-\frac{1}{2}\mathrm{tr}\ln(E^{-1}\sin E)} e^{-\frac{p^2\beta^2}{4s}} . \end{aligned} \qquad (18)$$

The first term is the zero temperature result [11], while the second term gives the finite temperature contribution as:

$$(\pi^2/45\beta^4) + O(\beta^2) ,$$

in the approximation that $\beta^2 B < 1$.

The full one loop effective action can now be written down in the two phases and yield the following expressions for the free energy densities*

* It is well known from the work of Coleman and Weinberg [12] that the phenomenon of dimensional transmutation occurs in the massless abelian Higgs model when the symmetry is broken by radiative corrections. This leads to the determination of λ in terms of e and the expectation value of φ. It is worth noting here that in the case of finite temperature, λ turns out to have a temperature dependence. This is seen from the following: Requiring

$$V = \frac{\lambda_R \varphi_R^4}{4!} + \frac{3e_R^4 \varphi_R^4}{64\pi^2}\left[\ln\frac{e_R^2\varphi_R^2}{\mu^2} - \ln 4\pi + \gamma - \frac{5}{6}\right] + \frac{e_R^2\varphi_R^2}{8\beta^2} - \frac{2\pi^2}{45\beta^4}$$

to have a minimum at a non-zero constant value φ_R^o yields the result ($e_R^2\varphi_R^{o2} = M_W^2$)

$$\lambda_R(\mu,\beta) = \frac{9e_R^4(\mu)}{8\pi^2}\left[-\ln\frac{M_W^2}{\mu^2} + \ln 4\pi - \gamma + \frac{1}{3} - \frac{4}{3\beta^2 M_W^2}\right] .$$

$$\mathcal{E}^{\text{broken}}(M_W,\beta) = -\frac{3M_W^4}{8(4\pi)^2} + \frac{(2\pi^2-1)}{(4\pi)^2}\frac{M_W^2}{\beta^2} - \frac{\pi^2}{45\beta^4} + O(\beta^2) \qquad (19)$$

and

$$\mathcal{E}^{\text{restored}}(B,\beta) = -\frac{1}{2}B^2 + \frac{1}{6}\frac{e^2B^2}{(4\pi)^2}\left[\ln\left(\frac{eB}{\mu^2}\right) + c + 1\right] - \frac{\pi^2}{45\beta^4} + O(\beta^2), \qquad (20)$$

where $c = -\ln 4\pi + \gamma - 1 + 12\zeta_R'(-1) = 0.0618$.

These are the finite temperature generalisations of the analogous expressions obtained by Shore [5] for $\beta = \infty$ or $T = 0$. Using these, we can obtain the critical B vs. T curve along which the two energy densities are equal at some finite T in the presence of an external magnetic field B along the z-axis. This leads to the result

$$-a_1 + a_2X^2 = -\frac{1}{2}Y^2 + \frac{e^2Y^2}{6(4\pi)^2}\left(\frac{1}{2}\ln e^2Y^2 + c + 1\right), \qquad (21)$$

where $a_1 = \frac{3}{8(4\pi)^2}$, $a_2 = \frac{(2\pi^2-1)}{(4\pi)^2}$,

and the dimensionless quantities X and Y are defined to be

$$X^2 = \frac{1}{\beta^2 M_W^2} , \qquad Y^2 = \frac{B^2}{M_W^4} ,$$

we have also chosen the convenient renormalisation point $\mu^2 = M_W^2$. This curve is plotted in Fig. 2 and we note that the critical field $B_c(0) = 0.07\, M_W^2$ and the critical temperature $KT_c(0) = 0.14\, M_W$.

Observe that, as expected, the critical field is indeed reduced at non-zero temperature from the zero temperature value. It is worthwhile noting that the constants a_1 and a_2 which determine the shape can be expected, in more realistic gauge models, to depend on the choice of gauge group and matter content of the theory and could give rise to different shapes for the critical curve which might allow restoration at drastically reduced external fields in the presence of small external temperatures. This is being investigated at present.

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FIGURE CAPTIONS

Fig. 1: External field B vs. temperature T plot.

Fig. 2: Plot of $x = \frac{kT}{M_w}$ versus $y = \frac{B}{M_w^2}$. The elliptical curve gives the critical points. The intercepts give the critical values of temperature (magnetic field) in the absence of magnetic field (temperature).

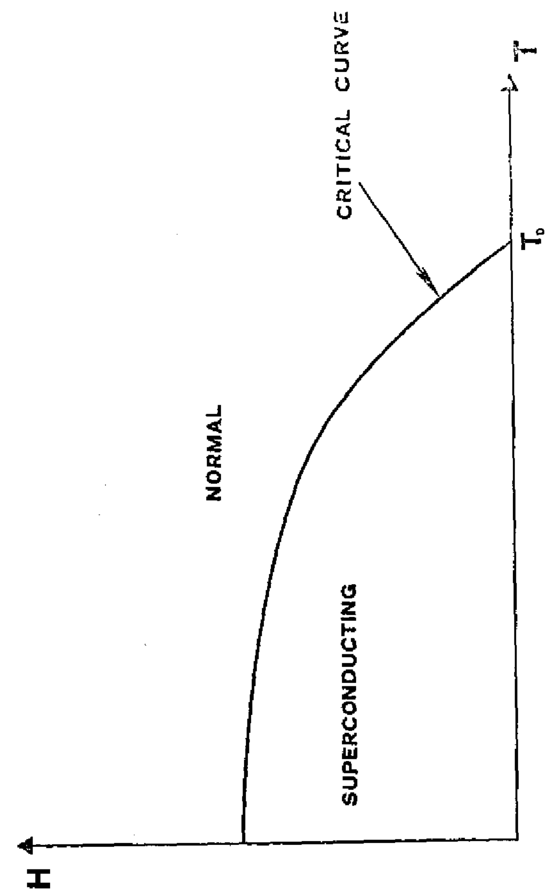

FIG 1

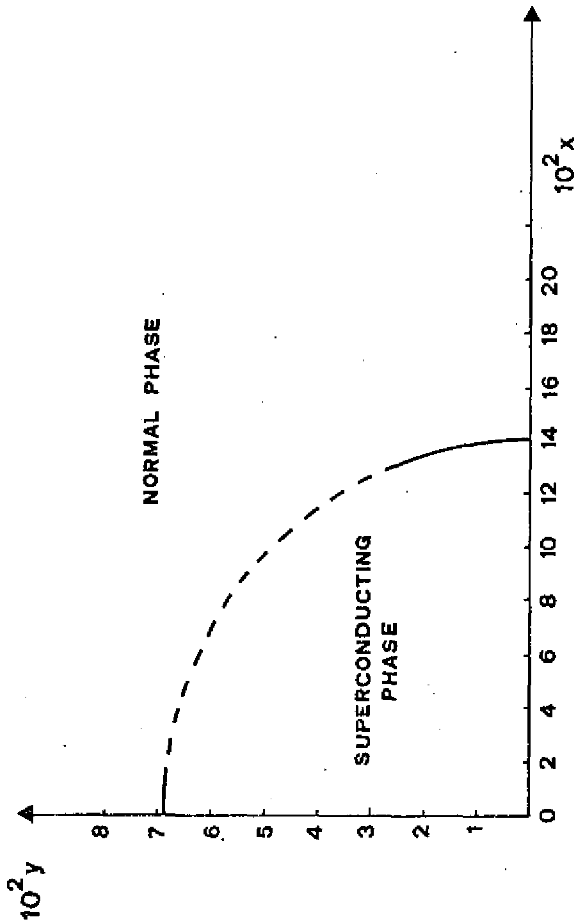

FIG 2

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