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ANOMALOUS PARAMAGNETIC BEHAVIOR: THE ROLE OF THE ZEROPOINT ELECTROMAGNETIC FLUCTUATIONS

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Abstract

The interaction of a magnetic dipole and the inductor of a RLC circuit without batteries, is described using the approach of stochastic electrodynamics. The purpose of this study is to clarify the effects of the current fluctuations on the paramagnetic behavior of a sample of magnetic material which is close to the solenoid. It is predicted a suppression in the average magnetization even in the case in which the circuit temperature is arbitrarily close to the absolute zero.

1. Introduction

Some years ago Boyer[1] has shown how an essentially classical model could account for the paramagnetic behaviour of a magnetic dipole in an external magnetic field. The classical model makes use of the random electromagnetic fluctuations associated with the zeropoint and the thermal radiations of Stochastic Electrodynamics (SED)[2]. After that, it was shown by Barranco, Brunini and França[3], that his model provides indeed a very good description of the observed paramagnetic behaviour at high and low temperatures.

More recently some authors[4–7] have addressed themselves to the study of the interaction between a microscopic subsystem, an electric dipole, for instance, and its environment, in an effort to clarify certain ideas brought to light by the Casimir effect[2]. We propose then to study the interaction between a rigid magnetic dipole, spinning around an external magnetic field, and a macroscopic solenoid which is part of a RLC circuit. Our aim is to show the environment induced modifications in the paramagnetic behaviour of the magnetic dipole and to understand the dynamical role played by the large energy reservoir associated with the stochastic zeropoint electromagnetic field.

In section 2 we review Boyer's model for free space paramagnetism in order to extend his method to treat, in section 3, the paramagnetic behaviour of a magnetic dipole in interaction with the inductor of a RLC circuit. In section 4 we resume our conclusions.

2. Free space paramagnetism

In this section we briefly review the treatment presented by Boyer[1] concerning the paramagnetic behaviour of a rigid magnetic dipole $\vec{\mu}$ placed at the origin of a coordinate system and subjected to an external constant magnetic field $\vec{B}_0 = B_0\ \hat{e}_z$ and to the fluctuating zeropoint plus thermal electromagnetic fields. In the dipole approximation the fluctuating magnetic field is given by

$$\vec{B}_{VF}(t) = \sum_{\lambda=1}^{2} \int d^3k\ \hat{\epsilon}(\vec{k},\lambda)\, H(\omega,T) \cos(\omega t + \xi(\vec{k},\lambda)) \quad , \tag{2.1}$$

where $\omega = c|\vec{k}\,|$ is the angular frequency of the plane wave of wave vector $\vec{k}$. The unit vectors $\hat{\epsilon}(\vec{k},\lambda)$ characterize the state of polarization of each plane wave and satisfy

$$\vec{k} \cdot \hat{\epsilon}(\vec{k},\lambda) = 0 \quad , \tag{2.2}$$

$$\sum_{\lambda=1}^{2} \hat{\epsilon}_i(\vec{k},\lambda)\hat{\epsilon}_j(\vec{k},\lambda) = \delta_{ij} - \frac{k_i k_j}{k^2} \quad , \tag{2.3}$$

and $\xi(\vec{k},\lambda)$ are random phases[1].

The function $H(\omega,T)$ is associated with the energy spectral density of both the zeropoint and the thermal fields at temperature T and it is given by[1]

$$H^2(\omega,T) = \frac{\hbar\omega}{2\pi^2} \coth\left(\frac{\hbar\omega}{2kT}\right) \quad . \tag{2.4}$$

The magnetic dipole $\vec{\mu}$ is related to the spin angular momentum vector $\vec{S}$ by

$$\vec{\mu} = \frac{eg}{2mc}\vec{S} \quad , \tag{2.5}$$

where e is the electric charge of a particle of mass m, and gyromagnetic factor g. The equation of motion of the spinning dipole is the Bhabha equation [1]

$$\dot{\vec{S}} = \vec{\mu} \times \vec{B}_0 + \vec{\mu} \times \vec{B}_{VF}(t) + \frac{2}{3c^3}\vec{\mu} \times \dddot{\vec{\mu}} \quad . \tag{2.6}$$

The first term in the right hand side corresponds to a precession around the z axis with frequency

$$\eta = \frac{\mu B_0}{S} \quad , \tag{2.7}$$

where $S = |\vec{S}\,|$ and $\mu = |\vec{\mu}\,|$. In this motion, precession without fluctuation and dissipation, the spin vector $\vec{S}$ makes a constant angle θ_0 relative to the z axis. The second term in (2.6) describes a random torque generated by the fluctuating magnetic field $\vec{B}_{VF}$ and the last term is the selftorque (due to the interaction between the magnetic dipole and its own radiation reaction field). These two last terms can be considered small in comparison with $\vec{\mu} \times \vec{B}_0$.

The net effect of the terms contained in the equation (2.6) is to produce an alignment probability distribution $P(\theta)$ associated with the angle of orientation θ between $\vec{\mu}$ and $\vec{B}_0$. Boyer[1] was able to find $P(\theta)$ by using a stationary Fokker-Planck equation for $P(\theta)$, namely

$$-P(\theta)\,\frac{\langle\Delta\theta\rangle}{\tau} + \frac{1}{2}\frac{d}{d\theta}\left[\frac{\langle(\Delta\theta)^2\rangle}{\tau}\,P(\theta)\right] = 0 \quad , \tag{2.8}$$

where $\langle\Delta\theta\rangle$ and $\langle(\Delta\theta)^2\rangle$ are the first and the second moments, respectively, for the change $\Delta\theta$ in θ during a small, but otherwise arbitrary, time τ.

In order to evaluate those moments Boyer used a perturbative quasi-Markovian approximation in equation (2.6) and found that

$$\frac{\langle\Delta\theta\rangle}{\tau} = \frac{2}{3c^3}\frac{\mu}{B_0}\eta^4 \sin\theta + \frac{2\pi^2}{3c^3}\left(\frac{\mu}{S}\right)^2 \cot\,\theta\,\eta^2\,H^2(\eta,T) \quad , \tag{2.9}$$

where the first term is due to radiation reaction and the last one is due to the random torque caused by the zeropoint and thermal fields. Boyer also found that[1]

$$\frac{\langle(\Delta\theta)^2\rangle}{\tau} = \frac{4\pi^2}{3c^3}\left(\frac{\mu}{S}\right)^2 \eta^2\ H^2(\eta,T) \quad . \tag{2.10}$$

Using (2.9) and (2.10) it is possible to find the exact solution of the Fokker-Planck equation (2.8) which is

$$P(\theta) = A\ \mathrm{sen}\,\theta\ \exp\left\{-\frac{S\eta}{\pi^2 H^2(\eta,T)}\ \cos\,\theta\right\} \quad , \tag{2.11}$$

where the constant A is determined by the normalization condition. Then the average component of the magnetic dipole $\vec{\mu}$ along the direction of the magnetic field $\vec{B}_0$, namely μ_z, can be calculated using the above probability distribution function, as follows:

$$\begin{aligned} \langle \mu_z \rangle &= \frac{eg}{2mc} \langle S_z \rangle \\ &= \frac{eg}{2mc} S \int_0^\pi d\theta \, P(\theta) \cos(\theta) \quad . \end{aligned} \tag{2.12}$$

Using (2.11) and (2.12) we get

$$\frac{\langle \mu_z \rangle}{g\mu_0} = \frac{S}{\hbar} \mathcal{L} \left\{ \frac{S\eta}{\pi^2 H^2(\eta, T)} \right\} \quad , \tag{2.13}$$

where $\mathcal{L}(x)$ is the Langevin function given by

$$\mathcal{L}(x) = \coth(x) - \frac{1}{x} \tag{2.14}$$

and μ_0 is the Bohr magneton, $\mu_0 = e\hbar/2mc$. Explicitly, equation (2.13) can be written as

$$\frac{\langle \mu_z \rangle}{g\mu_0} = \frac{S}{\hbar} \coth \left\{ \frac{2S}{\hbar \coth \left(\frac{\hbar\eta}{2kT} \right)} \right\} - \frac{1}{2} \coth \left(\frac{\hbar\eta}{2kT} \right) \quad . \tag{2.15}$$

Although the Boyer prediction given by equation (2.15) differs in form from the usual Quantum Mechanics prediction, that is,

$$\left(\frac{\langle \mu_z \rangle}{g\mu_0} \right)_{QM} = \left(j + \frac{1}{2} \right) \coth \left[(2j+1) \frac{\mu_0 B_0}{kT} \right] - \frac{1}{2} \coth \left(\frac{\mu_0 B_0}{kT} \right) \quad , \tag{2.16}$$

where $j(j+1) = \langle \vec{S}^2 \rangle_{QM} / \hbar^2$, it was shown by Barranco et al.[3] that, taking into account the experimental data, it is almost impossible to distinguish between (2.15) and (2.16).

3. Interaction between the magnetic dipole and the inductor

In this section we consider the modification of the paramagnetic behaviour of a spinning magnetic dipole $\vec{\mu}$ caused by the interaction with a solenoid which is part of a RLC circuit without batteries or any other source of power. The solenoid has radius a and N turns of wire uniformly distributed along its length $\ell (\ell \gg a)$. It is in the xy plane, in the $y < 0$ region, with its axis pointing parallelly to the x axis and at a distance $|y|$ from this (see Figure 1).

The equation that governs the fluctuating current $I(t)$ in the RLC circuit is such that[8]

$$L\dot{I}(t) + RI(t) + \frac{1}{C} \int dt \, I(t) = \mathcal{E}_N(t) + \mathcal{E}_{\text{dip}}(t) \quad , \tag{3.1}$$

where $\mathcal{E}_N(t)$ is the Nyquist electromotive force. The spectral distribution associated with $\mathcal{E}_N(t)$ has thermal and zeropoint contributions which are experimentally observed [9]. We are denoting by $\mathcal{E}_{\text{dip}}(t)$ the electromotive force generated in the circuit by the oscillating magnetic dipole which produces a time varying magnetic field through each coil of the solenoid. The calculation of this electromotive force is straightforward and the result is[10]

$$\mathcal{E}_{\text{dip}}(t) = \frac{N}{\ell} \frac{2\pi a^2}{cy^2} \dot{\mu}_x(t) \quad . \tag{3.2}$$

While discussing the modifications in the paramagnetic behaviour of the magnetic dipole we can safely neglect $\mathcal{E}_{\text{dip}}$, for it is very small when compared with $\mathcal{E}_N$ under the conditions we employ in this paper. For our purposes it is then legitimate to rewrite equation (3.1) approximately as

$$L\dot{I}_N(t) + RI_N(t) + \frac{1}{C} \int dt \, I_N(t) = \mathcal{E}_N(t) \quad , \tag{3.3}$$

$I_N(t)$ being the part of the current associated with the Nyquist fluctuations.

The variable electric current flowing along the coils of the solenoid will generate a magnetic field at the position of the dipole. This field is oriented along the x direction and it is given by[10]

$$\vec{B}_{\rm sol} = \hat{e}_x \frac{N}{\ell} \frac{2\pi a^2}{cy^2} I_N(t) \quad . \tag{3.4}$$

To account for any modification in the paramagnetic behaviour of the magnetic dipole we must include a new torque term, caused by the field of the solenoid at the position of the dipole, in the Bhabha equation (2.6), which now must read as

$$\dot{\vec{S}} = \vec{\mu} \times \vec{B}_0 + \vec{\mu} \times \vec{B}_{VF} + \vec{\mu} \times \vec{B}_{\rm sol} + \frac{2}{3c^3} \vec{\mu} \times \dddot{\vec{\mu}} \quad , \tag{3.5}$$

where $\vec{B}_{\rm sol}$ is given by (3.4). The extra term in equation (3.5) alters the alignment probability distribution discussed in the previous section. As the new torque is very small, when compared with the deterministic torque, we are allowed to use the perturbative quasi-Markovian approach again in order to evaluate the first moment $\langle \Delta\theta \rangle / \tau$ and the second moment $\langle (\Delta\theta)^2 \rangle / \tau$ which appear in equation (2.8).

To achieve this goal it is convenient to use for the Nyquist current a Fourier decomposition that is similar in form to that used by Boyer[1] for describing the statistical properties of $\vec{B}_{VF}(t)$. We shall denote

$$I_N(t) = \int_0^\infty d\omega \, J(\omega, T) \cos(\omega t + \zeta(\omega)) \quad , \tag{3.6}$$

where $\zeta(\omega)$ is an element of a set of random phases, uniformly distributed over the range $[0, 2\pi]$, and statistically independent of the set of random phases associated with the magnetic field $\vec{B}_{VF}(t)$. The function $J(\omega, T)$ depends on the frequency ω and the circuit temperature T and will be given below.

Performing calculations similar to those of Boyer we find that

$$\frac{\langle \Delta\theta \rangle}{\tau} = \frac{2}{3c^3} \frac{\mu}{B_0} \eta^4 \sin\theta + \frac{2\pi^2}{3c^3} \left(\frac{\mu}{S}\right)^2 \eta^2 H^2(\eta, T_d) \cot\theta + $$
$$+ \frac{\pi}{8} \left(\frac{N}{\ell} \frac{2\pi a^2}{cy^2}\right)^2 \left(\frac{\mu}{S}\right)^2 J^2(\eta, T_c) \cot\theta \quad , \tag{3.7}$$

and that

$$\frac{\langle (\Delta\theta)^2 \rangle}{\tau} = \frac{4\pi^2}{3c^3} \left(\frac{\mu}{S}\right)^2 \eta^2 H^2(\eta, T_d) + \frac{\pi}{4} \left(\frac{N}{\ell} \frac{2\pi a^2}{cy^2}\right)^2 \left(\frac{\mu}{S}\right)^2 J^2(\eta, T_c) \quad . \tag{3.8}$$

In these equations we have explicitly indicated the possibility that the dipole be in a thermal bath at temperature T_d while the circuit is in another thermal bath at a temperature T_c *different* from T_d.

To these new results corresponds a stationary Fokker-Planck equation (in the form (2.8)) the exact solution of which is

$$P(\theta) = \text{const.} \sin\theta \exp\left\{ \frac{-S\eta \cos\theta}{\pi^2 H^2(\eta, T_d) + \frac{3\pi c^3}{16\eta^2} \left(\frac{N}{\ell} \frac{2\pi a^2}{cy^2}\right)^2 J^2(\eta, T_c)} \right\} \quad . \tag{3.9}$$

The average values of the component of the magnetic moment $\vec{\mu}$ along the z axis is then given by a Langevin function of a different argument, namely

$$\langle \mu_z \rangle = \frac{egS}{2mc} \mathcal{L} \left\{ \frac{S\eta}{\pi^2 H^2(\eta, T_d) + \frac{3\pi c^3}{4\eta^2} \left(\frac{N}{\ell} \frac{2\pi a^2}{cy^2}\right)^2 \frac{\langle |\tilde{\mathcal{E}}_N(\eta, T_c)|^2 \rangle}{|\mathcal{Z}(\eta)|^2}} \right\} \quad . \tag{3.10}$$

Here we used the relation[8] that $J^2(\eta, T_c)$ bears with the spectral density of the mean squared current, i.e.

$$J^2(\omega, T) = 4 \frac{\langle |\tilde{\mathcal{E}}_N(\omega, T)|^2 \rangle}{|\mathcal{Z}(\omega)|^2} \quad , \tag{3.11}$$

where $\tilde{\mathcal{E}}_N(\omega,T)$ is the Fourier transform of the Nyquist electromotive force and $\mathcal{Z}(\omega)$ is the impedance of the RLC circuit,

$$\mathcal{Z}(\omega) = R - i\left(\omega L - \frac{1}{\omega C}\right) \quad . \tag{3.12}$$

In order to fully appreciate the result given in equation (3.10) it is important to notice that [8,9],

$$\langle |\tilde{\mathcal{E}}_N(\omega)|^2 \rangle = \frac{R\hbar\omega}{2\pi} \coth\left(\frac{\hbar\omega}{2kT_c}\right) \quad . \tag{3.13}$$

Thus equation (3.10) can be put in the form

$$\begin{aligned} \frac{\langle \mu_z \rangle}{g\mu_0} &= \frac{S}{\hbar} \coth\left\{ \frac{2S/\hbar}{\coth\left(\frac{\hbar\eta}{2kT_d}\right) + \left(\frac{a}{y}\right)^4 \in(\eta) \coth\left(\frac{\hbar\eta}{2kT_c}\right)} \right\} - \\ &- \frac{1}{2}\left\{ \coth\left(\frac{\hbar\eta}{2kT_d}\right) + \left(\frac{a}{y}\right)^4 \in(\eta) \coth\left(\frac{\hbar\eta}{2kT_c}\right) \right\} \quad , \end{aligned} \tag{3.14}$$

where

$$\begin{aligned} \in(\eta) &= \frac{3\pi^2 N^2 cR}{\ell^2 \eta^2 |\mathcal{Z}(\eta)|^2} \\ &= \frac{(3\pi^2 N^2 c/\ell^2 \eta^2 R)}{1 + \frac{L^2}{R^2}\frac{1}{\eta^2}\left(\eta^2 - \frac{1}{LC}\right)^2} \quad . \end{aligned} \tag{3.15}$$

4. Presentation of the results and discussion

Comparing equations (2.15) and (3.14) we see that the effect of the presence of the solenoid in the neighborhood of the dipole amounts to a correction in the observed paramagnetic behaviour of the dipole that depends on the parameters characterizing the RLC circuit, and especially the solenoid with which the dipole interacts, as well as on the angular frequency η characteristic of the precession of the dipole. The presence

of the solenoid is capable of altering sensitively the statistical behaviour of the magnetic dipole as we will see in a moment.

To illustrate the effects of the inductor on the magnetic dipole we shall present a plot of equation (3.14) illustrating the variation of $\langle \mu_z \rangle / \mu_0$, as a function of B_0, for an ion like Gd^{3+} that has a ratio $S/\hbar$ (the SED quantity equivalent[3] to $\sqrt{\langle \vec{S}^2 \rangle / \hbar^2}$ in Q.M.) equal to 4. The solenoid employed in the calculation is similar to those used by Möllenstedt and Bayh[11] to observe the Aharonov-Bohm effect. To be more specific we have chosen the following set of parameters: $N/\ell = 2 \times 10^3 \mathrm{cm}^{-1}$, $a = 7 \times 10^{-4}\mathrm{cm}$, $|y| = 1.4 \times 10^{-3}\mathrm{cm}$, $R = 10^{-10}\mathrm{sec/cm}$, $L = 5 \times 10^{-20}\ \mathrm{sec}^2/\mathrm{cm}$, and $C = 10^3\mathrm{cm}$. The same parameters were used recently by Blanco et al.[6] and Dechoum and França[7] in their study of the interaction between the solenoid and an electric dipole.

Figure 2 shows the magnetization (in units of the Bohr magneton μ_0) as a function of the applied magnetic field B_0 for a dipole at $T_d = 2K$. The upper curve represents the free-space paramagnetic behaviour of the dipole as predicted by equation (2.15) and experimentally confirmed[3]. The two lower curves were calculated from equation (3.14) in the cases $T_c = 2K$ and $T_c = 300K$. We see that in both cases the magnetization is strongly affected by the solenoid. Observe also that the higher the temperature of the circuit the lower the magnetization for each value of B_0. This is expected on physical grounds.

To stress the importance of the zeropoint fluctuating field to the result (3.14) we made another plot of the magnetization. This time we compare the free-space magnetization with the magnetization we would obtain in the presence of a solenoid at

$T_c = 0K$. According to (3.14), we get

$$\frac{\langle \mu_z \rangle}{g\mu_0} = \frac{S}{\hbar} \coth \left\{ \frac{2S/\hbar}{\coth\left(\frac{\hbar\eta}{2kT_d}\right) + \left(\frac{a}{y}\right)^4 \in(\eta)} \right\}$$

$$-\frac{1}{2} \left\{ \coth\left(\frac{\hbar\eta}{2kT_d}\right) + \left(\frac{a}{y}\right)^4 \in(\eta) \right\} , \tag{4.1}$$

for $\coth\left(\frac{\hbar\eta}{2kT_c}\right)$ tends to unity when T_c approaches zero. Figure 2 shows all the resulting curves for different temperatures of the circuit ($T_c = 0, 2$ and $300K$) when all the other parameters remain unchanged. The important role played by zeropoint fluctuations, existing in the RLC circuit and generating the modification of the paramagnetic behaviour of the dipole, may be viewed by noting that at $B_0 = 10KG$ the ratio $\frac{\langle \mu_z \rangle_{\text{free-space}} - \langle \mu_z \rangle}{\langle \mu_z \rangle_{\text{free-space}}}$ is 0.67. We also note that the zeropoint current fluctuations are responsible for 87% of the modification in $\langle \mu_z \rangle$ in the case $T_c = 2K$.

These results may be interpreted as follows. The RLC circuit picks up thermal and zeropoint energy from the environment and then the solenoid radiates part of this energy to the dipole. Thus the dipole wins its tendency to alignment along the direction of the applied field $\vec{B}_0$. In turn, the dipole also radiates back to the RLC circuit a part of the energy that it picks up from the environment, as we shall show in a forthcoming paper. The net result is an increase in the energy of the dipole which affects its average orientation angle.

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Figure Captions

Figure 1

Schematic picture showing the magnetic dipole $\vec{\mu}$ and the solenoid orientation with respect to the applied magnetic field $\vec{B}_0$ and the coordinate system. The fluctuating current $I(t)$ is also indicated.

Figure 2

Magnetization per ion as predicted by our equations (3.14) and (4.1). The ratio $|y|/a = 2$ and the values of the parameters of the RLC circuit are given in section 4. The dotted curve is associated with the normal paramagnetic behavior (eq. (2.15) with $T = 2K$), and the experimental points are indicated by the small triangles. The other curves show the anomalous magnetization due to presence of the solenoid. The absolute temperature of the circuit is indicated by T_c. The paramagnetic sample is maintained at $T_d = 2K$.

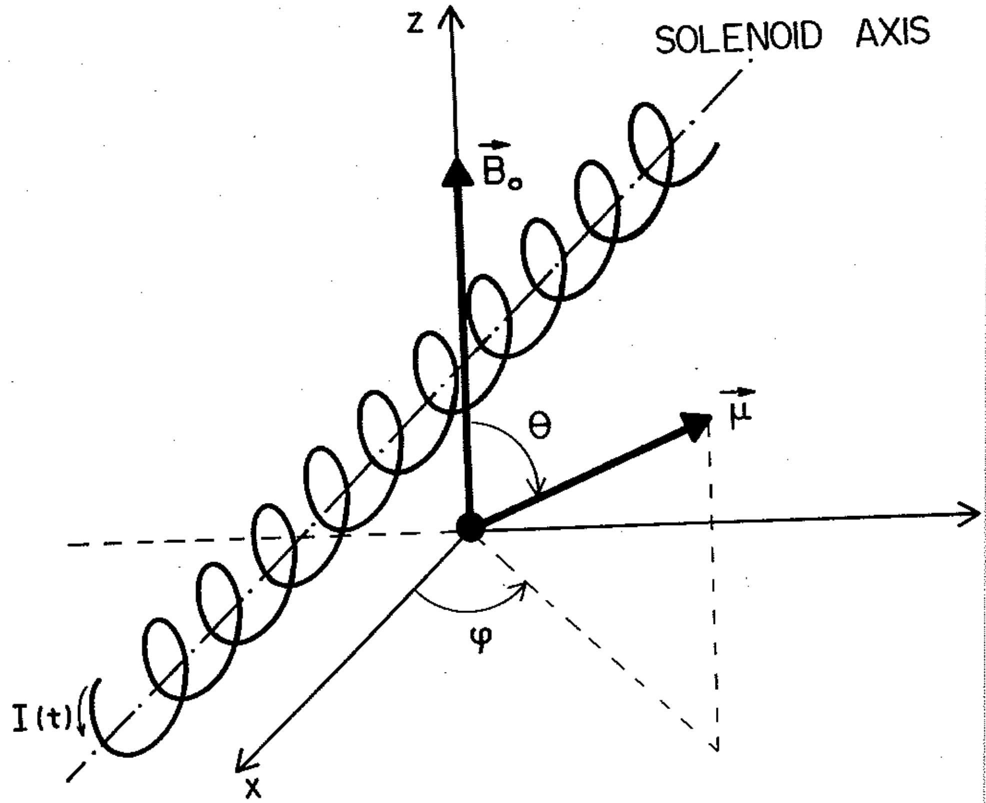
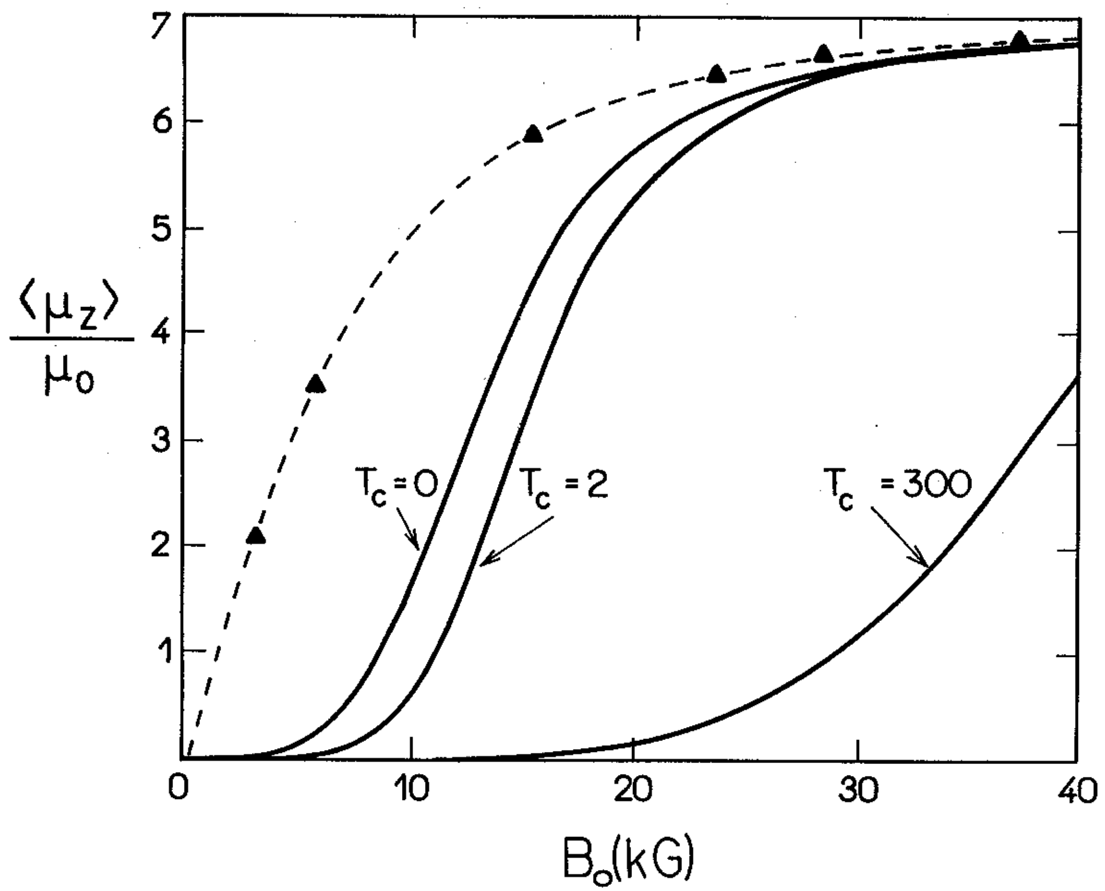


FIGURE 1



$\frac{\langle \mu_z \rangle}{\mu_0}$
0
1
2
3
4
5
6
7
10
20
30
40
$T_c = 0$
$T_c = 2$
$T_c = 300$
B_o(kG)

FIGURE 2