IC/83/113
INTERNAL REPORT
(Limited distribution)

International Atomic Energy Agency

and

United Nations Educational Scientific and Cultural Organization

INTERNATIONAL CENTRE FOR THEORETICAL PHYSICS

THEORETICAL DESCRIPTION OF THE PROPERTIES OF MAGNETIZATION FLUCTUATIONS IN THE VICINITY OF PHASE TRANSITION FROM PARAMAGNETIC PHASE TO FERROMAGNETIC PHASE WITH DOMAIN STRUCTURE *

W. Wasilewski

International Centre for Theoretical Physics, Trieste, Italy.

MIRAMARE - TRIESTE

August 1983

* Not to be submitted for publication.

** Permanent address: Department of Physics, Technical University, Malczewskiego 29, 26-600 Radom, Poland.

1. Introduction

In real finite-size ferromagnetic samples in the ferromagnetic phase the ground state is a homogeneous magnetization state or inhomogeneous magnetization state, i.e. domain structure [1,2]. The kind of the domain structure is determined by a solid geometry of a sample, by a type of anisotropy energy and a temperature of a sample. In further considerations we confine ourselves to examine magnetic materials with uniaxial anisotropy. We will also assume that the anisotropy energy of a material is less than the maximal value of the demagnetizing energy. We consider a flat-parallel film made of uniaxial ferromagnet with the easy magnetization axis perpendicular to the surfaces of the sample and parallel to the z axis of the cartesian coordinate system. We assume that the dimensions of the sample in the (x,y) plane are much larger than the thickness D along the z axis, hence in further considerations they will be treated as infinite. It is well known [1,2] that in this kind of a sample the state of homogeneous magnetization in the plane of the film (i.e. in hard magnetic direction) is the ground state when the thickness D of the film is less than the critical thickness D_c. For $D = D_c$ there is a continuous phase transition from the homogeneous magnetization state F_{HM} to the domain structure F_{DS}. Domain structure is the ground state when the thickness D of the film is greater than D_c. The essential physical reason for the existence of the domain structure is the demagnetizing field determined by the overall dimensions and solid geometry of the sample. In the case of flat-parallel film which we investigate the thickness D decides which of the competitive energies, anisotropy or demagnetizing, is dominant and define the ground state. Both of the ferromagnetic ground states i.e. F_{HM} and F_{DS} can be realized for the range of temperature $T < T_c^*(D)$. The symbol $T_c^*(D)$ denotes the temperature of continuous phase transition from a paramagnetic phase P to the ferromagnetic phase F_{HM} for $D \leqslant D_c$ or F_{DS} for $D > D_c$. It is illustrated on the phase diagram (Fig. 1).

We present the theoretical description of the phase transition $P \Rightarrow F_{DS}$ induced by the temperature change for the case $D \gg D_c$ within the phenomenological theory of phase transitions.

2. The statical properties of magnetization fluctuations.

We represent the free energy F of a sample in the form:

$$F = F_i\{|\bar{M}|\} + F_s\{\bar{M}, \nabla\bar{M}\} \quad . \qquad (1)$$

The isotropic part $F_i\{|\bar{M}|\}$ of the free energy describes strong exchange interactions which decide the average thermodynamical value of the magnetization vector $M_o(T) = |\langle\bar{M}(\bar{r})\rangle_T|$. The term $F_s\{\bar{M}, \nabla\bar{M}\}$ describes the free energy connected with weak long-distance interactions. The magnetization fluctuations $\bar{m}(\bar{r})$ are defined as: $\bar{m}(\bar{r}) = \bar{M}(\bar{r}) - \langle\bar{M}\rangle_T$. Near the phase transition temperature the term $F_1\{M\}$ of the free energy can be [3] written as:

$$F_1\{M\} = \frac{1}{2}\int_V \left(a\tau M^2 + b M^4 \right) d\vec{r} \qquad (2)$$

where a, b are the standard material constants appearing in the Landau form (2) of the free energy. The symbol τ denotes the reduced temperature: $\tau = (T - T_c)/T_c$ and T_c is the Curie temperature of the infinite isotropic ferromagnet. In the case of a flat-parallel sample made of a uniaxial ferromagnet with the easy magnetization axis perpendicular to the surfaces of the sample (and parallel to the z axis of the cartesian coordinate system) the term $F_2\{\vec{M}, \nabla\vec{M}\}$ can be represented [4] in the form:

$$F_2\{\vec{M}, \nabla\vec{M}\} = \frac{1}{2}\int_V \left[\alpha(\nabla\vec{M})^2 - \beta M_z^2 - \vec{h}^d \vec{M} \right] d\vec{r} \qquad (3)$$

where α is the macroscopic exchange constant, $\beta < 4\pi$ is the constant of uniaxial anisotropy. The symbol $\vec{h}^d = \vec{H}^d(\vec{r})$ denotes the vector of the demagnetizing field. In the paramagnetic phase $(M_o(T > T_c) = 0)$ the free energy of fluctuations $\vec{m}(\vec{r})$ is described by the functional:

$$F\{\vec{m}, \nabla\vec{m}\} = F\{\vec{M}, \nabla\vec{M}\} - F\{\vec{M}_o, \nabla\vec{M}_o\} =$$

$$= \frac{1}{2}\int_V \left[a\tau m^2 + b m^4 + \alpha(\nabla\vec{m})^2 - \beta m_z^2 - \vec{h}^d \vec{m} \right] d\vec{r}. \qquad (4)$$

The demagnetizing field $\vec{h}^d$ can be easily expressed in terms of $\vec{m}$ using the Maxwell equations in magnetostatic approximation, i.e.

$$\mathrm{rot}\,\vec{h}^d = 0 \quad ; \quad \mathrm{div}\left(\vec{h}^d + 4\pi\vec{m}\right) = 0 . \qquad (5)$$

Hence, for Fourier transforms $\vec{m}(\vec{q})$ and $\vec{h}^d(\vec{q})$ of the quantities $\vec{m}(\vec{r})$, $\vec{h}^d(\vec{r})$ we obtain the following relation:

$$\vec{h}^d(\vec{q}) = -4\pi q^{-2}(\vec{q}\,\vec{m})\vec{q} \qquad . \qquad (6)$$

In a bilinear approximation with respect to $\vec{m}(\vec{q})$ the fluctuations free energy ΔF can be written in the form:

$$\Delta F = \sum_{\vec{q}} \mathcal{E}_{ij}(\vec{q},\tau)\, m_i(\vec{q})\, m_j(-\vec{q}) \quad ; \quad i,j = z,\perp . \qquad (7)$$

The elements of the matrix $\mathcal{E}_{ij}(\vec{q},\tau)$ in (7) are given by:

$$\mathcal{E}_{zz}(\vec{q},\tau) = a\tau + \alpha q^2 + 4\pi q^{-2}k^2 - \beta ,$$
$$\mathcal{E}_{\perp\perp}(\vec{q},\tau) = a\tau + \alpha q^2 + 4\pi q^{-2}\varkappa^2 , \qquad (8)$$
$$\mathcal{E}_{z\perp}(\vec{q}) = \mathcal{E}_{\perp z}(\vec{q}) = 4\pi q^{-2} k\varkappa ,$$

where the following notation has been used:

$$\vec{q} = (\varkappa_x, \varkappa_y, k) \; ; \; \vec{\varkappa} = (\varkappa_x, \varkappa_y, 0) \; ; \; q^2 = \varkappa^2 + k^2 .$$

The magnetization fluctuations $\vec{m}(\vec{r})$ should fulfill appropriate boundary conditions on surfaces of the sample. In the case of strong surface anisotropy ("pinning") these boundary conditions take the form:

$$m_z(z=0) = m_z(z=D) = 0 \quad ; \quad \left.\frac{\partial m_\perp}{\partial z}\right|_{z=0} = \left.\frac{\partial m_\perp}{\partial z}\right|_{z=D} = 0 . \qquad (9)$$

From the boundary conditions (9) we obtain:

$$k \to k_n = n\pi D^{-1} \quad ; \quad n = 1,2,\ldots \qquad (10)$$

As usual, the stability condition for the postulated paramagnetic state is the positive definiteness of the bilinear form (7) describing the fluctuation energy ΔF. In our case it means that the determinant $\Delta(\bar{q},\tau)$ of the matrix $\mathcal{E}_{ij}(\bar{q},\tau)$ should be positive:

$$\Delta(\bar{q}_n,\tau) = \left|\mathcal{E}_{ij}(\bar{q}_n,\tau)\right| = (\alpha q_n^2 + a\tau)^2 + (4\pi-\beta)(\alpha q_n^2 + a\tau) - 4\pi\beta q_n^{-2} k_n^2 \geqslant 0, \qquad (11)$$

where $\bar{q}_n = (\varkappa_x, \varkappa_y, k_n)$. From the analysis of expression (11) it follows [5] that, for fixed values of D, n, τ the function $\Delta(\bar{q}_n,\tau)$ has the minimum for the value $q_{0,n}$ of q_n given by formulae:

$$q_{0,n}^2(D,\tau) = n q_{0,1}^2(D,0)\left[1 - \frac{a\tau}{4\pi-\beta}\right],$$
$$q_{0,1}^2(D,0) = \sqrt{\frac{4\pi\beta}{(4\pi-\beta)\alpha}}\;\frac{\pi}{D} \qquad (12)$$

Inserting (12) into (11) we obtain an expression describing the minimal value of the determinent $\Delta(\bar{q}_n,\tau)$. From the expression (11) it follows also that $\Delta(\bar{q}_{0,n},\tau) > \Delta(\bar{q}_{0,1},\tau)$ for $n > 1$. Hence, we can limit ourselves to the case $n = 1$ in our stability examination. The dependence $\Delta(\bar{q}_1,T)$ is illustrated in Fig. 2. The minimal value $\Delta_0(T) = \Delta(\bar{q}_{0,1},T)$ is a function only temperature. The function $\Delta_0(T)$ changes its sign for $T = T_c^*(D)$:

$$\Delta_0(T>T_c^*) > 0 \; ; \; \Delta_0(T=T_c^*) = 0 \; ; \; \Delta_0(T<T_c^*) < 0 . \qquad (13)$$

For the temperature $T < T_c^*(D)$ the paramagnetic state is no longer a stable one. From the formulae (11) - (13) we obtain the following expression for $T_c^*(D)$:

$$T_c^*(D) = T_c\left\{1 + a^{-1}\left[\beta - 4\sqrt{\frac{\pi\alpha\beta}{4\pi-\beta}}\;\frac{\pi}{D}\right]\right\}. \qquad (14)$$

The temperature $T_c^*(D)$ given by the formula (14) is [6] the temperature of the phase transition from the paramagnetic phase to the ferromagnetic phase with domain structure of a period $\lambda(D,T) = 2\pi\varkappa_0^{-1}$ in (x,y) plane. For the range of temperature $T \leqslant T_c^*$ we have obtained [5] the following formula describing the dependence of the period λ of the domain structure on the thickness D of the film and temperature:

$$\lambda(D, T\leqslant T_c^*) = \lambda(D,T_c^*) + A(D)\left[1 - \frac{T}{T_c^*}\right], \qquad (15)$$

where the initial period $\lambda(D,T_c^*)$ is given by the formula:

$$\lambda(D,T_c^*) = \frac{2D}{\sqrt{2DD_c^{-1}-1}} \quad ; \quad D_c = \sqrt{\pi\alpha\beta^{-1}(4\pi-\beta)}\,. \qquad (16)$$

The dependence of parameter $A(D)$ in (15) on thickness D is given by:

$$A(D) = A_1 D + A_{1/2} D^{1/2} \qquad (17)$$

where

$$A_1 = 2a D_c^{-1} \quad ; \quad A_{1/2} = \frac{3\sqrt{2}\,a}{4\pi-\beta} D_c^{-1/2}.$$

It has been shown above that in ferromagnetic phase the period of the domain structure depends on temperature, and near the point of phase transition this dependence is linear. Thus, it is in agreement with results of experimental research [7,8] carried out on the orthoferrites samples on the basis of Faraday's effect near the Neel temperature. The obtained experimental data are presented in Fig. 3.

Obtained theoretical results described above are consistent with experimental results presented in [12] for $T \geq T_c$ and in [7,8] for temperatures $T \lesssim T_c$. In this context conclusions formulated by Drabkin et al. [14] on the basis of experimental result are not clear. In the above mentioned paper the results of ferromagnetic materials properties measurements in critical region obtained by polarized neutron diffraction method were presented. Drabkin et al. deduced the following conclusions from obtained experimental results. In Curie temperature quasistable phase of homogeneous magnetization appear. Below T_c in the temperature $T_d < T_c$ quasistable phase of homogeneous magnetization changes to the phase with domain structure. The region $T_c - T_d$ where the phase of homogeneous magnetization exist can be from ~ 0.1 K to ~ 3 K what depends on the kind of a material.

3. Correlation functions of magnetization fluctuations in paramagnetic phase

Taking the formulae (12), (14) into account we can rewrite the formula (11) describing the determinant $\Delta(\bar{q}_n, \tau)$ in the form:

$$\Delta(\bar{q}_n, \tau) \Rightarrow \Delta(\bar{q}_n, \tau_n^*) = (4\pi - \beta)\left[a\tau_n^* + \alpha q_{o,n}^2 \left(\frac{q_n}{q_{o,n}} - \frac{q_{o,n}}{q_n}\right)^2\right], \quad (18)$$

where the following notation has been used:

$$\tau_n^* = \frac{T - T_{c,n}^*}{T_{c,n}^*} \quad ; T_{c,n}^* = T_c^*(D) - 2a^{-1}(n-1)T_c \alpha q_{o,1}^2 . \quad (19)$$

The probability $P(\bar{q}_n)$ of the appearence of fluctuations $\bar{m}(\bar{q}_n)$ is given [9] by the formula:

$$P(\bar{q}_n) \sim \exp\left[-\frac{1}{k_B T} \varepsilon_{ij}(\bar{q}_n, \tau_n^*) m_i(\bar{q}_n) m_j(-\bar{q}_n)\right] \quad (20)$$

The correlation functions can be calculated [9] from the shape of a probability distribution (20) and from the formulae (8) describing the elements of matrix $\varepsilon_{ij}(\bar{q}_n, \tau)$. The static correlation function in $\bar{q}$ - representation have the form:

$$\langle m_z(\bar{q}_n) m_z(-\bar{q}_n)\rangle \sim \frac{\varepsilon_{11}(\bar{q}_n, \tau_n^*)}{\Delta(\bar{q}_n, \tau_n^*)} ,$$
$$\langle m_\perp(\bar{q}_n) m_\perp(-\bar{q}_n)\rangle \sim \frac{\varepsilon_{zz}(\bar{q}_n, \tau_n^*)}{\Delta(\bar{q}_n, \tau_n^*)} , \quad (21)$$
$$\langle m_z(\bar{q}_n) m_\perp(-\bar{q}_n)\rangle \sim \frac{\varepsilon_{z1}(\bar{q}_n)}{\Delta(\bar{q}_n, \tau_n^*)} .$$

In the $\vec{r}$ - representation the correlation functions are given by formulae:

$$G_{zz}(\vec{r},\vec{r}') \sim \sum_n \int d\vec{\varkappa} \langle m_z(\vec{q}_n) m_z(-\vec{q}_n)\rangle \sin k_n z \sin k_n z' \exp[i\vec{\varkappa}(\vec{r}_\perp - \vec{r}_\perp')],$$

$$G_{\perp\perp}(\vec{r},\vec{r}') \sim \sum_n \int d\vec{\varkappa} \langle m_\perp(\vec{q}_n) m_\perp(-\vec{q}_n)\rangle \cos k_n z \cos k_n z' \exp[i\vec{\varkappa}(\vec{r}_\perp - \vec{r}_\perp')], \qquad (22)$$

$$G_{z\perp}(\vec{r},\vec{r}') \sim \sum_n \int d\vec{\varkappa} \langle m_z(\vec{q}_n) m_\perp(-\vec{q}_n)\rangle \sin k_n z \cos k_n z' \exp[i\vec{\varkappa}(\vec{r}_\perp - \vec{r}_\perp')],$$

where: $\vec{r}_\perp = (x, y, 0)$; $z, z' \in (0, D)$.

Any of the correlation functions (21) , (22) are singular for $\vec{q}_n = \vec{q}_{0,n}$ and simultaneously $\tau = \tau_n^* = 0$. The phase transition from the paramagnetic phase to ferromagnetic one with domain structure takes place for $\vec{q} = \vec{q}_{0,1}$ and $T = T_c^*(D)$. It means that full information on the domain structure appearing at the temperature T_c^* contains the behavior of critical fluctuations in paramagnetic phase. Thus, we can conclude that even in a paramagnetic phase "material knows" which domain structure will be produced at the transition point.

4. The dynamical properties of magnetization fluctuations.

The form of the equations of motion describing the magnetization fluctuations $\vec{m}(\vec{r},t)$ in the paramagnetic phase can be determined [9] under the assumption that the state of the system is fully described at a moment t by the vector $\vec{m}(\vec{r},t)$ and by the derivative $\partial\vec{m}/\partial t$. Let us also assume, what follows from the experimentally discovered fact [10] of the existence of spin waves in the paramagnetic phase: that the equations of

which describe small oscillations of the magnetization fluctuation vector $\vec{m}(\vec{r},t)$ around the equilibrium position $\vec{m}(\vec{r},t) = 0$. In this case the kinetic energy K of the oscillations of the magnetization fluctuation vector is the functional which can be written in the form:

$$K = \frac{1}{2}\int_V \left(\Gamma \frac{\partial\vec{m}}{\partial t}\right)^2 d\vec{r}\,. \qquad (23)$$

The parameter Γ in (23) is treated as a macroscopic material constant. The Lagrangian $L\{\vec{m}, \nabla\vec{m}, \partial\vec{m}/\partial t\}$ describing the fluctuations $\vec{m}(\vec{r},t)$ is defined [11] as $L = K - \Delta F$. Taking the formulae (4) and (23) into consideration we obtain:

$$L = \frac{1}{2}\int_V \left\{ \left(\Gamma\frac{\partial\vec{m}}{\partial t}\right)^2 - \alpha(\nabla\vec{m})^2 - a\tau\vec{m}^2 - b\vec{m}^4 + \beta m_z^2 + \vec{h}^{d}\vec{m} \right\} d\vec{r}\,. \qquad (24)$$

In description of the dynamics of magnetization fluctuations $\vec{m}(\vec{r},t)$ we must take into account energy dissipation processes. Within the phenomenological theory which we use , those processes are described by means of the functional R of the form:

$$R = \frac{1}{2}\int_V \varrho\left\{\frac{\partial\vec{m}}{\partial t}\right\} d\vec{r} \qquad (25)$$

with the density $\varrho\left\{\frac{\partial\vec{m}}{\partial t}\right\}$ given by:

$$\varrho\left\{\frac{\partial\vec{m}}{\partial t}\right\} = \sum_i \sum_j \Lambda_{ij} \frac{\partial m_i}{\partial t}\frac{\partial m_j}{\partial t} \; ; \; i,j = z,\perp\,. \qquad (26)$$

The symbols $\Lambda_{ij} = \Lambda_{ji}$ denote kinetic coefficients. The equations of motion for the magnetization fluctuations have, in our case [11] , the form:

$$\frac{\delta L}{\delta m_i} = - \frac{\partial \varsigma}{\partial \left(\frac{\partial m_i}{\partial t} \right)} \tag{27}$$

It is easy to see that for $\Gamma = 0$ the equations of motion (27) pass to the standard Onsager equations of motion:

$$\frac{\delta (\Delta F)}{\delta m_i} = \Lambda_{ij} \frac{\partial m_j}{\partial t} , \tag{28}$$

which describe the process of relaxation of the magnetization fluctuations $\bar{m}(\bar{r},t)$. The demagnetizing field $\bar{h}^d$ is connected with the magnetization vector $\bar{m}$ via Maxwell equations (5). Moreover the magnetization fluctuation vector $\bar{m}(\bar{r},t)$ should fulfill the boundary conditions (9). In further considerations we adopt the simplifying assumption $\Lambda_{ij} = \delta_{ij}\Lambda$. Let us introduce the Fourier transforms $\bar{m}(\bar{q},\omega)$ of the $\bar{m}(\bar{r},t)$:

$$m_i(\bar{r},t) = \int d\bar{q} \int d\omega \, m_i(\bar{q},\omega) \exp(i\omega t + i\bar{q}\bar{r}) . \tag{29}$$

Substituting (28) into the equations of motion (26) and set of Maxwell equation (5) with formulae (24), (26) and the boundary conditions (9) taken into account we can obtain the dispersion relations of the spin waves in the paramagnetic phase. Near the Curie temperature there are two modes of spin-wave-like oscillations: the optic mode of spin waves and the soft mode of spin waves. In the process of phase transition from the paramagnetic phase to ferromagnetic phase with domain structure a decisive role is played by the soft mode of spin waves. The dispersion relation of the soft mode has the form:

$$\omega = i \frac{\Lambda}{2\Gamma^2} \pm \left[\omega_s^2(\bar{q}_n, \tau_n^*) - \frac{\Lambda^2}{4\Gamma^4} \right]^{1/2} \tag{30}$$

where $\omega_s(\bar{q}_n,\tau_n^*)$ is the frequency of undamped oscillations of the soft mode. It is given by formula:

$$\omega_s(\bar{q}_n,\tau_n^*) = \Gamma^{-1} \left[\varepsilon_{zz}(\bar{q}_n,\tau_n^*) + \varepsilon_{\perp\perp}(\bar{q}_n,\tau_n^*) \right]^{-1/2} \sqrt{\Delta(\bar{q}_n,\tau_n^*)} \tag{31}$$

The symbols τ_n^*, $\varepsilon_{ij}(\bar{q}_n, \tau_n^*)$ were defined above. The symbol $\Delta(\bar{q}_n, \tau_n^*)$ denotes the determinant of matrix ε_{ij} (see formula (18)). As it has been shown $\lim_{T \to T_c^*} \Delta(\bar{q}_{0,1},T) = 0$ and $\Delta(q_{0,1},T < T_c) < 0$. It means that in the vicinity of phase transition temperature the formula (30) describes the overdamped oscillations when the inequality $\omega_s(\bar{q}_n, \tau_n) < \Lambda/2\Gamma^2$ is fulfilled. This effect in vicinity of phase transition temperature has been experimentally observed [10]. In such a range of T and q_1 where $\omega_s(q_1,T) \ll \Lambda/2\Gamma^2$ we obtain from (30) and (31) the formulae describing the roots of equation (30):

$$\omega^{(1)} = i\Lambda\Gamma^{-2} ,$$

$$\omega^{(2)} = i\tilde{T}^{-1}(q_n,\tau_n^*) = i\Lambda^{-1}\Gamma^2\omega_s^2(\bar{q}_n,\tau_n^*); \quad \tilde{T}^{-1} \sim \Delta(\bar{q}_n,\tau_n^*). \tag{32}$$

The dependence on time of fluctuations of magnetization describes, in this case, the processes of relaxation :

$$m_i(t) \sim \exp(-\Lambda\Gamma^{-2}t) + \sum_n m_i(\bar{q}_n,\omega^{(2)}) \sin k_n z \; \times \tag{33}$$

$$\exp(i\bar{\varkappa}\bar{r}_\perp) \exp\left(-\tilde{T}^{-1}(\bar{q}_n,\tau_n^*)t\right).$$

The symbol $\tilde{T} = \tilde{T}(\bar{q}_n, \tau^*_n)$ denotes the relaxation time. From the preceding study of the properties of the function $\Delta(\bar{q}_n, \tau^*_n)$ it follows that the fluctuations which relax most slowly are those described by $\bar{q}_n = \bar{q}_{o,1}$ i.e. those which generate the phase transition from the paramagnetic phase to the ferromagnetic phase with domain structure in (x,y) plane of the period:

$$\lambda(D,\tau_c^*) = 2\pi \varkappa_0^{-1}(D,\tau_c^*) = 2\pi \left[q^2_{o,1}(D,\tau_c^*) - \pi^2 D^{-2} \right]^{-1/2} \qquad (34)$$

From the properties of determinant $\Delta(\bar{q}_{o,1}, \tau^*)$ we obtain:

$$\lim_{T \to T_c^*} \tilde{T}(\bar{q}_{o,1}, T) = \infty \qquad (35)$$

The expression (35) describes the effect of a critical slowing-down of the fluctuations generating the phase transition. On the other hand, it was shown above that the probability of appearance of fluctuations described by the value $\bar{q} = \bar{q}_{o,1}$ is the greatest. In our opinion these fluctuations have been observed by Rekveldt, van Woesik and Meijer [12] as a "... kind of domain structure in the paramagnetic region..." .

5. Temperature shift of the main maximum of neutron critical scattering.

From the first measurements of neutron magnetic critical scattering till the experiment performed by Stump and Maier [13] it was believed that the intensity of neutron critical scattering had always its maximum at the temperature equal to the Curie temperature T_c. Stump and Maier found for Ni that the temperature $T_m > T_c$ for which the maximum of scattering appears depends on the observation angle Θ for a fixed neutron wavelength λ, i.e. depends on scattering vector $\bar{Q}$. The scattering vector $\bar{Q}$ is determined as follows: $\bar{Q} = \bar{k}_f - \bar{k}_i$ where the symbol $\bar{k}_f$ denotes the wave vector of scattered neutron and $\bar{k}_i$ is the wave vector of incoming neutron. The result obtained in [13] means that the temperature shift $\Delta T = T_m - T_c$ is the function of the scattering vector:

$$\Delta T = T_m - T_c = f(Q) \qquad (36)$$

Bally et al. [15] have observed the temperature shift in Co and Fe and found that in some region of value of the scattering vector Q the function $f(Q)$ has the form:

$$f(Q) \sim Q^3 \qquad \text{for Co},$$

$$f(Q) \sim Q^4 \qquad \text{for Fe}.$$

Experiments performed by Blinowski and Ciszewski [16] showed that the temperature shift ΔT depends not only on the observation angle Θ but also on the wavelength λ_i of incoming neutrons, what is illustrated in Fig. 4. The effect of temperature shift has been described in many papers (see review article [17]) on the basis of Van Hove [18] theory of critical scattering.

We propose a different physical mechanism which causes the temperature shift of the maximum of neutron critical scattering. Let us assume that experimentally observed temperature shift is a result of scattering of neutrons on the spin waves in paramagnetic phase. Let us also assume that in the process of neutron scattering the following laws of conservation are fulfilled:

$$\mathcal{E}(Q) = \hbar\,\omega_s(q,T) \quad ; \quad \bar{Q} + \bar{q} = 0 \qquad (37)$$

The energy transfer $\mathcal{E}(Q)$ and the value Q for the case of small-angle neutron scattering are determined as:

$$\mathcal{E}(Q) = \frac{\hbar^2}{2m}\left|k_f^2 - k_i^2\right| \approx \frac{\hbar^2}{2m}Q^2 \quad ; \quad Q = \left|\bar{k}_f - \bar{k}_i\right| \approx k_i\,\theta \,. \qquad (38)$$

where m is the neutron mass. The equations of motion (27) in linear approximation give the following dispersion relation for spin waves in paramagnetic phase of bulk material:

$$\omega_s(q,\tau) = \Gamma^{-1}\sqrt{a\tau + \alpha q^2} \quad ; \quad \tau = \frac{T - T_c}{T_c} \,. \qquad (39)$$

The term describing the demagnetizing energy and the dissipation term have been neglected in the equations of motion. Substituting (38), (39) into (37) we obtain:

$$\Delta T = T_m - T_c = B\,Q^2\left(Q^2 - Q_o^2\right) \qquad (40)$$

The symbol Q_o denotes the threshold value of scattering vector. As long as Q is less than Q_o no shift of temperature should occur. The threshold value Q_c is described as follows:

$$Q_c = \frac{2m}{\Gamma\hbar}\sqrt{\alpha} \,. \qquad (41)$$

The quantity B in (40) is determined as follows:

$$B = \left(\frac{\Gamma\hbar}{2m}\right)^2 \frac{T_c}{a} \,. \qquad (42)$$

As it has been mentioned above the results illustrated in Fig. 4. have been obtained with the help of small-angle neutron critical scattering technique. In this case we can assume that:

$$Q = k_i\,\theta = \frac{2\pi}{\lambda_i}\,\theta \,. \qquad (43)$$

Substituting (43) into (40) we obtain:

$$T_m - T_c = C(\lambda_i)\,\theta^2\left[\theta^2 - \theta_c^2(\lambda_i)\right] \qquad (44)$$

where

$$C(\lambda_i) = \left(\frac{\pi\hbar\Gamma}{m}\right)^2 \frac{T_c}{a}\,\lambda_i^{-4} \,. \qquad (45)$$

The threshold value $\theta_c(\lambda_i)$ of the observation angle is given by formula:

$$\theta_c(\lambda_i) = \frac{2m\sqrt{\alpha}}{h\Gamma}\,\lambda_i \,. \qquad (46)$$

The value of macroscopic exchange constant α for iron is equal to: $\alpha \approx 1.4\times10^{-16}\,m^2$. The Curie temperature of examined [16] sample is equal to: $T_c = 1043$ K. From the experimental data illustrated in Fig. 4. we can estimate the values:

$\frac{\Delta\theta}{\Delta\lambda} \approx$ [illegible] rad/m , $C(\lambda_i) \approx 3.5\times10^{7}\,K/rad^4$.

So from formulae (45), (46) we obtain:

$$\Gamma = \frac{2m\sqrt{x}}{\hbar}\left(\frac{\partial \varepsilon}{\partial \lambda}\right)^{-1} \approx 3.6\times 10^{-10}\,s\,, \qquad (47)$$

$$a = \left(\frac{\hbar \,[illegible]}{m\lambda_c^2}\right)^2 \frac{T_c}{\varepsilon(\lambda_c)} \approx 1.6\times 10^{-[illegible]}. \qquad (48)$$

The physical model of critical scattering of neutrons presented above can be treated as the additional test of possibility of description of spin-wave-like excitations of magnetization fluctuations in paramagnetic phase within the phenomenological theory. On the other hand, the proposed physical mechanism causing the effect of temperature shift is very simple and quite different from all others used to explain this effect.

ACKNOWLEDGMENTS

The author would like to thank Professor Abdus Salam, the International Atomic Energy Agency and UNESCO for hospitality at the International Centre for Theoretical Physics, Trieste.

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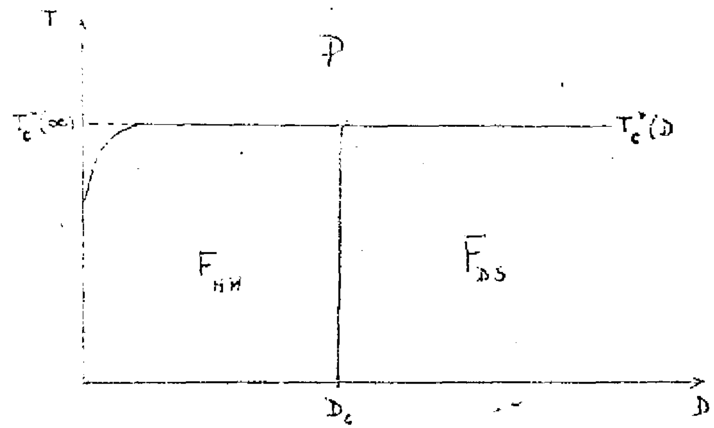

Fig. 1.

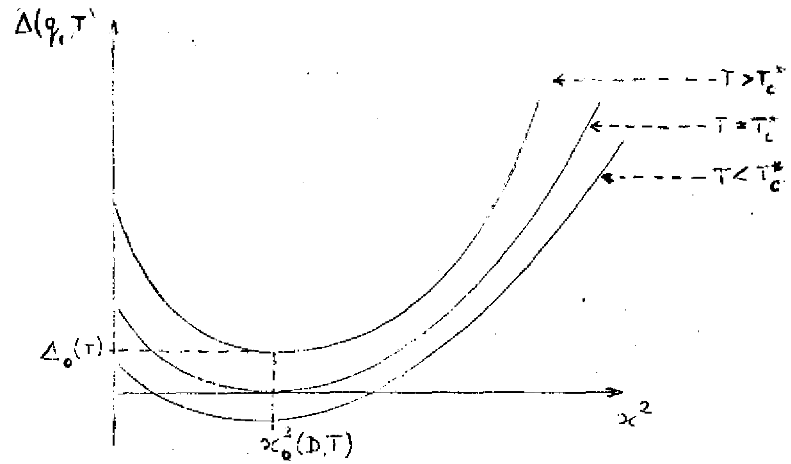

Fig 2

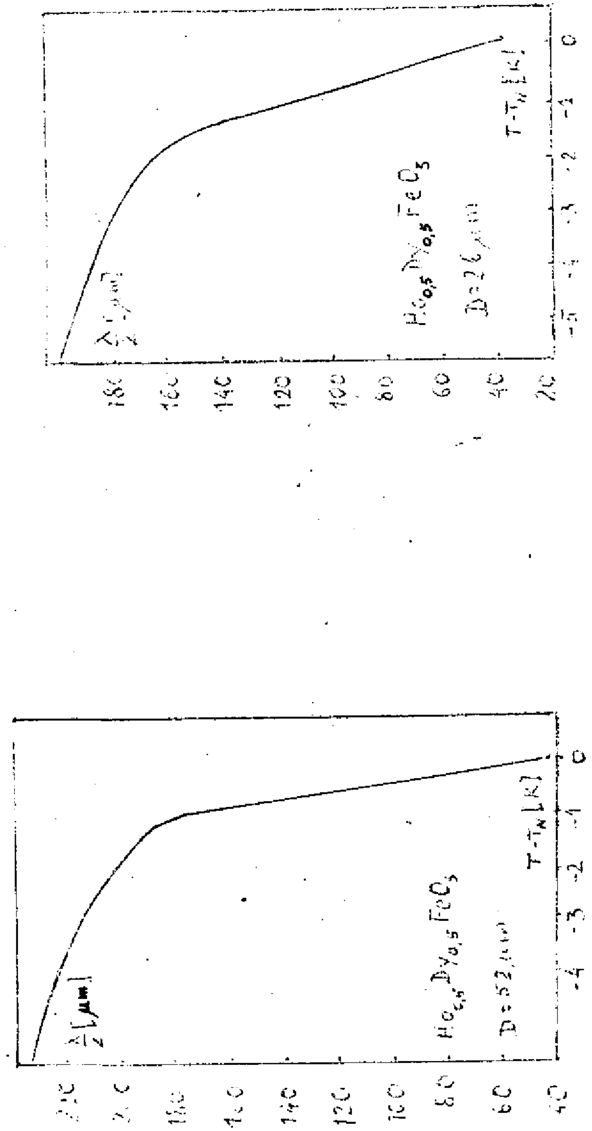

Fig. 3.

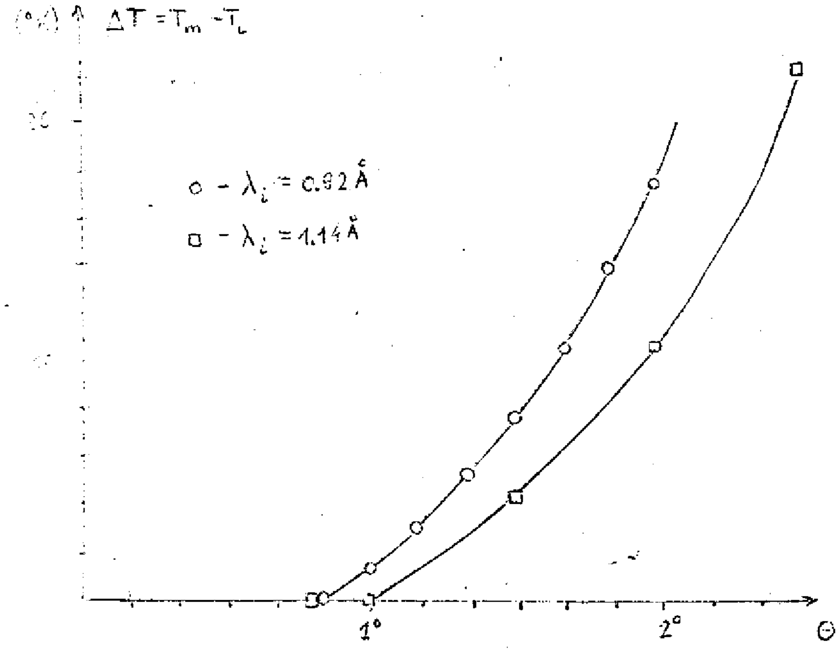

Fig. 4.

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