Parametric Excitation of Spin Waves in Antiferromagnetic $MnCO_3$

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Parametric excitation of spin waves of the low-frequency spectral branch in the weakly ferromagnetic substance $MnCO_3$ is observed when the high-frequency field and the static magnetic field applied to the basal plane of the crystal are parallel. The experiments were performed at frequencies $\nu_p = 25$–47 GHz and at liquid-helium temperature. The dependence of the threshold field strength on the temperature and on the wave vector is studied in detail and it is found that in the range from 1.2 to 2.1°K the spin-wave decay obeys the law $\Delta\nu_{1k}(k\to 0) = cT^{0.4}$. Two sharp peaks are observed at $T \leq 1.5$°K on the $h_c(k)$ curve. The linear dependence of $\nu(k)$ for one of them permits the conclusion that this peak is due to resonance interaction between excited spin waves and phonons. A value $\alpha = 0.79\times 10^{-5}$ kOe-cm is obtained for the exchange constant in the spin-wave spectrum on the basis of the known speed of sound in $MnCO_3$.

1. INTRODUCTION

In antiferromagnets with anisotropy of the "easy plane" type, parametric excitation of spin waves corresponding to the low-frequency branch of the spectrum is possible under definite conditions. This phenomenon was observed experimentally in[1-3] in the hexagonal antiferromagnet $CsMnF_3$. We have investigated parametric excitation of spin waves in the weak antiferromagnet $MnCO_3$[4]. The present paper is devoted to a detailed description of experimental excitation and study of spin-wave relaxation in $MnCO_3$.

The spin-wave spectrum for antiferromagnets with "easy plane" anisotropy was calculated by Borovik-Romanov[5] and Turov[6]. It consists of two branches described by the following formulas:

$$(\nu_{1k}/\gamma)^2 = H\sin\theta(H\sin\theta + H_D) + H_\Delta^2 + \alpha^2 k^2, \quad (1)$$

$$(\nu_{2k}/\gamma)^2 = 2H_A H_E + (H_D + H\sin\theta)H_D + H^2\cos^2\theta + \alpha^2 k^2, \quad (2)$$

where γ is the gyromagnetic ratio, H_D is the Dzyaloshinskiĭ field, H_E is the exchange field, H_Δ^2 is the gap due to the hyperfine interaction, H_A is the anisotropy field, α is the exchange constant, k is the wave number of the spin waves, and θ is the angle between the field H and the principal axis of the crystal. The spin oscillations corresponding to the first and second branches of the spectrum are described by different components of the vectors

$$\mathbf{l} = \mathbf{l}_0 + \lambda e^{i(\omega t - \mathbf{kr})}, \quad \mathbf{m} = \mathbf{m}_0 + \mu e^{i(\omega t - \mathbf{kr})}$$

and are excited by different components of the microwave field h.

As shown by Ozhogin[7], the anharmonicity in such systems is determined by the existence of a nonlinear connection between the branches of the spectrum. From the solution of the nonlinearized system of equations of motion it follows that at large amplitudes h at the frequency ν_1, the motion of the moments is described by the complete assembly of the components of the vectors λ and μ, and the components that determine in the linear approximation the oscillations of the second branch oscillate with frequency $2\nu_1$. This leads to a possibility of parametric excitation of the first branch of the spectrum by the electromagnetic wave of frequency $2\nu_1$, the polarization of which corresponds to excitation of the second branch ($\mathbf{h} \parallel \mathbf{H}$), and the amplitude exceeds the "threshold" value h_c determined by the dissipation in the system. The parametric excitation of the spin waves is revealed by the appearance of absorption in the sample in fields H corresponding to excitation of the lower branch at the frequency ν_1 with $k \geq 0$.

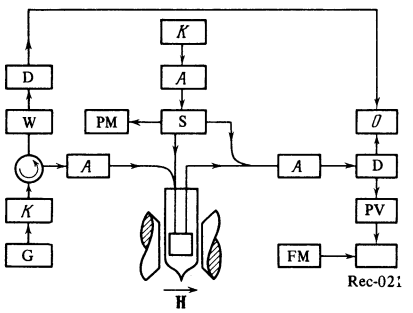

FIG. 1. Block diagram of spectrometer: A–attenuator, S–switch, D–detector, K–klystron, W–wave meter, PV–peak-value voltmeter, O–oscilloscope, FM–field meter, PM–power meter, G–sawtooth pulse generator, REC–automatic recorder.

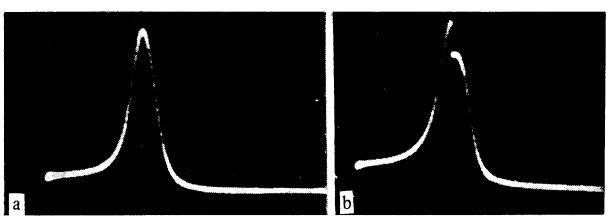

FIG. 2. Oscillogram of signal transmitted to resonator at $H < H_0$: a) $h < h_c$, b) $h > h_c$.

2. PROCEDURE AND SAMPLES

A block diagram of the spectrometer is shown in Fig. 1. The source of the microwave power was a klystron oscillator operating in the 8 mm band, and its frequency was swept slowly. The microwave signal was applied through a circulator and a calibrated attenuator to a cylindrical copper resonator operating in the H_{012} mode with $Q \approx 10000$, in which the investigated sample was placed. The detected signal, proportional to the power P_1 passing through the resonator, was fed to the vertical plates of an oscilloscope. The oscilloscope

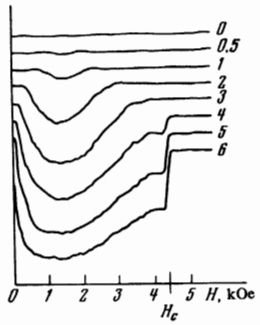

FIG. 3. Automatically plotted amplitudes of signal transmitted to resonator as function of the static field H. The numbers on the curves show the value of the microwave power in the resonator in decibels. T = 1.7°K, ν_p = 36.1 GHz.

screen displayed the frequency characteristic $P_1(\nu_p)$ of the resonator. When the threshold absorption took place in the sample, the amplitude $P_1 max$ and the shape of this curve changed (see Fig. 2). The dependence of $P_1 max$ on the static magnetic field was measured with a peak-value voltmeter and registered with an x-ray recording potentiometer (see Fig. 3).

To determine the amplitude h of the microwave field, we measured the power fed to the resonator, its Q, and the coupling coefficient β. The instantaneous value of the microwave power was measured with a square-law detector whose sensitivity was systematically calibrated during the course of the experiment. The calibration was with the aid of a second klystron operating in the pulsed regime at the resonator frequency. The pulse power was measured with a standard thermistor instrument. The measurement accuracy was 10%.

The field h at the sample was calculated from the known formulas for the H_{012}-mode field distribution. It was assumed that the sample does not change the field configuration in the resonator significantly, since the tested substance had $\mu \approx 1$ and the influence of ϵ was small because the sample was placed in a node of the electric field. The absolute accuracy with which h was obtained was 15% and was determined mainly by the errors in the measurement of Q (10%), of the coupling coefficient β (10%), and of the power P (10%). The change of h in each series of experiments (at fixed Q and β) can be measured more accurately (5%).

To determine the value of the critical field $\tilde{h}_c$ we measured the power P_c at which threshold absorption set in. The value P_c determined from the oscillograms and from the potentiometer chart coincided within 0.1 dB. It was observed in the experiment that $\tilde{h}_c$ depends on the time of action of the microwave field on the antiferromagnet. We have therefore performed the experiments at a constant rate of frequency deviation 2GHz/sec. In this case $\tilde{h}_c$ differed from the value h_c measured in the static regime by not more than 5%.

The measurements were performed at pump frequencies ν_p = 25–47 GHz.

In many experiments it was necessary, simultaneously with the excitation of the electron spin waves, to vary the spectrum gap H_Δ^2 due to the hyperfine interaction. This was done by saturating the nuclear subsystem of the Mn^{++} ions at the nuclear-resonance frequency ~ 640 MHz[8]. To this end, two turns of wire of 0.1 mm diameter were wound around the sample placed in the resonator, and this loop was coupled with the oscillator by a coaxial cable. The oscillator power (~1 W) was

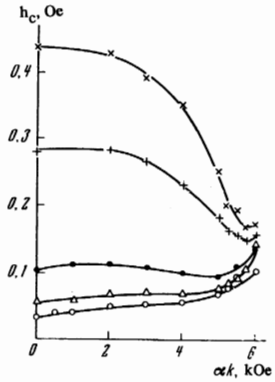

FIG. 4

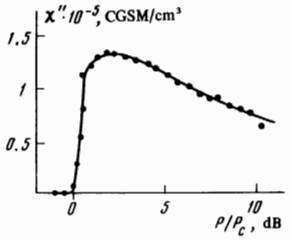

FIG. 5

FIG. 4. Dependence of the threshold field on the value of αk at different temperatures: X–2.17, +–2.03, ●–1.75, Δ–1.56, O–1.23°K; ν_p = 36.1 GHz.

FIG. 5. Dependence of above-threshold susceptibility χ'' on the power at αk = 3.5 kOe, ν_p = 36.3 GHz, and T = 1.2°K.

sufficient to obtain a nuclear-subsystem temperature T_n = 20°K; H_Δ^2 was decreased thereby to 0.5 kOe^2.

The static magnetic field was produced with a laboratory electromagnet or with a superconducting solenoid. The field-measurement accuracy was not worse than 1%. The measurements were performed in the temperature interval 2.1–1.2°K. To prevent possible overheating of the samples, the resonator was filled with superfluid helium. The bath temperature was measured with a resistance thermometer. The accuracy with which the temperature was measured and maintained was 0.01°K. We used in the experiments $MnCO_3$ single crystals grown by the hydrothermal method[1)].

3. EXPERIMENTAL RESULTS

Figure 3 shows a plot of the absorption curves at different amplitudes of the microwave field h at the sample. In these measurements, the sample was placed in the resonator in such a way that the magnetic fields H and h were parallel to each other and were in the basal plane of the crystal. It is seen from Fig. 3 that absorption in fields $H < H_0$ sets in at a certain value of the amplitude h_c. The value of the field H_0 coincides with that calculated by formula (1) for k = 0 and $\nu = \nu_p/2$. Thus, it can be assumed that the onset of absorption at $h > h_c$ is due to parametric excitation of the first branch of the spin-wave spectrum at the frequency $\nu_p/2$.

Figure 4 shows the experimentally determined dependence of h_c on αk at different temperatures at the frequency 36.1 GHz. The value of αk was calculated for the given H and T using formulas (1) with the constants taken from[8]. We note that at $\alpha k > 5.4$ kOe the sample is situated in fields $H < 1.5$ kOe. Measurements of the susceptibility[5] have shown that the sample is not saturated.

From the power ΔP absorbed by the sample beyond

[1)]The authors are grateful to M. Yu. Ikornikov and V. R. Gakel' for supplying the $MnCO_3$ crystals grown by them.

the spin-wave excitation threshold we can calculate the imaginary part of the above-threshold susceptibility χ'' using the formula

$$\chi'' = \Delta P / \nu_p \int h^2 dV.$$

The above-threshold susceptibility increases with increasing pump amplitude and passes through a maximum at $P/P_c \approx 3$ dB. Figure 5 shows a plot of $\chi''(P/P_c)$ for $\alpha k = 3.5$ kOe, $T = 1.2°K$ and $\nu_p = 36.3$ GHz.

At temperatures below 1.5°K, we have observed the appearance of singularities on the $P_{1max}(H)$ curves. Figure 6 shows the signal transmitted through the resonator at $T = 1.45°K$ and $\nu_p = 36.3$ GHz. In the fields H_1 and H_2, the absorption amplitude has a minimum, which can be regarded as an increase of ~0.2 dB in the threshold field h_c at these points. Such a small change of the threshold lies beyond the accuracy of the power measurement, but is reliably observed on the curves of the above-threshold absorption. In individual crystals, the peak in the field H_1 has an amplitude ~4 dB and nine satellites, with amplitudes 20–30 times smaller than the amplitude of the main peak, are observed on its two sides.

To establish the nature of these peaks it was necessary to ascertain whether the position of these peaks depends on the wave vector **k** or on the field **H**. This can be done by exciting, at the given frequency, a spin wave with another wave vector **k** in the field **H**. As follows from (1), it is necessary for this purpose to change the spectrum gap H_Δ^2 due to the hyperfine interaction. H_Δ^2 is inversely proportional to the temperature of the Mn^{++} nuclear system[9]. It was shown in[9] that the temperature of the nuclei can be raised to an arbitrary value T_m above that of the lattice by applying enough power at the nuclear-resonance frequency corresponding to the given temperature. When varying the temperature of the nuclei, we observed the following: when the nuclear system was saturated, the fields H_0 and H_2 increased and the field H_1 remained unchanged. In addition, the value of h_c for all values of k increased somewhat (within 0.1 dB), which may correspond to a sample temperature rise $\Delta T \approx 0.01°K$. From the change of H_0 we can calculate H_Δ^2 for the attained nuclear temperature T_n, and then calculate αk for the new value of H_2, It turned out that in the described experiment the field H_2 changed in such a way that the corresponding value of αk remained constant. We can therefore conclude that the anomaly in the field H_1 is determined by the value of the static

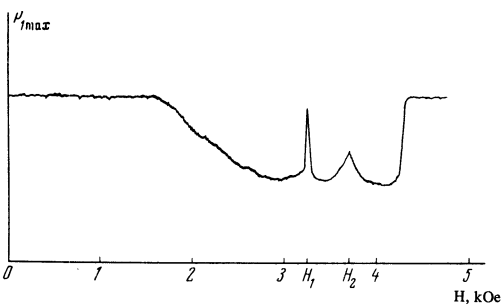

FIG. 6. Automatic plot of above-threshold absorption at $T = 1.45°K$ and $\nu_p = 36.3$ GHz.

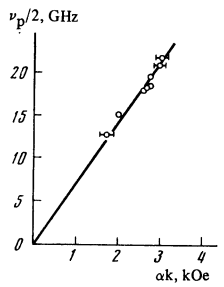

FIG. 7

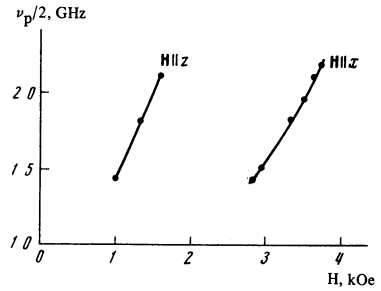

FIG. 8

FIG. 7. Dispersion law $\nu(\alpha k)$ for the anomaly in the field H_2.
FIG. 8. $\nu(H)$ dependence for the anomaly in the field H_1.

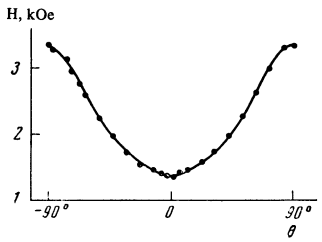

FIG. 9. Dependence of field H_1 on the angle θ between the direction of the static field and the principal axis of the crystal.

field **H**, and the anomaly in the field H_2 is determined by the wave vector **k**.

A series of experiments performed at the different pump frequencies has shown that H_1 and H_2 vary with frequency and do not depend on the orientation of the static field **H** in the basal plane. Figure 7 shows the results of these experiments for the peak in the field H_2, plotted in the coordinates $\nu_p/2$ and αk. The experimental points lie on a straight line passing through the origin:

$$\nu = {}^1/_2\nu_p = 7.1\,[\text{GHz/kOe}]\,\alpha k\,[\text{kOe}].$$

A plot of $\gamma(H)$ for the peak in the field H_1 is shown in Fig. 8.

Parametric excitation of spin waves was observed at different values of the angle θ between the static field **H** and the basal plane of the sample. It turned out that for a fixed value of αk the threshold field h_c remains constant within 10% when θ varies from 90 to 20°. In addition, it was established that when θ is varied the fields H_0 and H_2 vary in accordance with formula (1) at a fixed value of k.

Figure 9 shows the dependence of the field H_1 on the angle θ.

4. DISCUSSION OF RESULTS

1. The threshold field h_c, which is determined by the spin-wave relaxation frequency $\Delta\nu_k$, was measured in experiments on parametric excitation of spin waves. The question of parametric excitation of spin waves in rhombohedral antiferromagnets with weak ferromagnetism was considered by Ozhogin[7] and later by Kolganov[10]. For the case $\nu_p \ll \nu_2$ they obtained the following formula for the threshold field:

$$h_c = \Delta\nu_k \nu_p / \gamma^2 (2H_\perp + H_D), \qquad (3)$$

where $H_\perp$ is the field applied in basal plane. This formula differs significantly from the formula for anti-

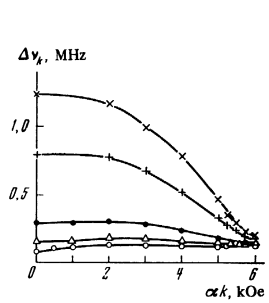

FIG. 10

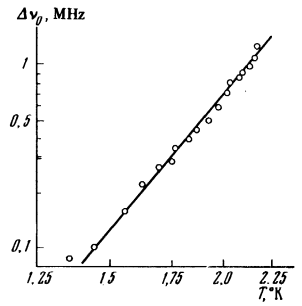

FIG. 11

FIG. 10. Dependence of the spin-wave damping $\Delta\nu_k$ on αk for different temperatures: X–T = 2.17, +–2.03, O–1.75, Δ–1.56, O–1.23°K.

FIG. 11. Temperature dependence of the spin-wave damping $\Delta\nu_0$ ($\alpha k \to 0$) (in logarithmic scale).

ferromagnets without weak ferromagnetism (e.g., $CsMnF_3$[1]) in that the denominator contains the sum $2H_\perp + H_D$. Accordingly, we have observed parametric excitation of spin waves in very weak fields (down to 10 Oe). Our experiments have also confirmed that, in accordance with (3), h_c is determined by the projection $H \sin\theta$ of the static field on the basal plane.

To determine experimentally the dependence of h_c on H and ν_p, it is necessary to obtain from other experiments data on the dependence of $\Delta\nu_k$ on the wave vector and on the magnetic field. At present there are no such data.

We have subsequently used the theoretical formula (3) to determine $\Delta\nu_k$ from the measured values of h_c.

2. Figure 10 shows the dependence of $\Delta\nu_k$ on the product αk at different temperatures. Attention is called to the decrease of $\Delta\nu_k$ with increasing wave vector. A similar $\Delta\nu_k(k)$ dependence was observed experimentally for $CsMnF_3$[2] and follows apparently from experiments on parametric excitation of a pair of electronic and nuclear spin waves and $RbMnF_3$ in fields stronger than the flipping field[11]. On the other hand, the field h_c in $CuCl_2 \cdot 2H_2O$[12] increases quadratically with increasing k. A similar dependence is usually observed in ferrites. We can therefore conclude that different spin-wave relaxation mechanisms can predominate in antiferromagnets with different magnetic structures. Another possible cause of such an appreciable difference in the $\Delta\nu_k(k)$ dependence might be the influence of the hyperfine interaction, which is many times stronger for Mn^{++} than for Cu^{++}. To determine the influence of the nuclear system on the spin-wave relaxation, the above-described experiments on saturation of nuclear system were performed and revealed that heating the nuclei to $T_n \approx 20°K$ leaves the magnitude and character of the $\Delta\nu_k(k)$ dependence practically unchanged.

The question of the relaxation time in ferromagnets with anisotropy of the easy plane type was considered theoretically by Ozhogin[13] and by Woolsey and White[14]. Ozhogin considered a three-magnon relaxation process. From the energy and the momentum conservation laws it follows that this process should be accompanied by production of short-wave magnons of the second branch of the spectrum. At the temperatures employed by us, the energy of such magnons is $h\nu_{2k} \gg k_BT$, which leads to an exponential dependence of the relaxation rate on the temperature. In our case this corresponds to an increase of $\Delta\nu_k$ by three orders of magnitude when the temperature is changed from 1.2 to 2.2°K, something not confirmed by the experiment, although at T = 2.1°K the experimental value of $\Delta\nu_k$ coincides with that calculated from the formulas of[13]. In the opinion of Woolsey and White[14], the principal relaxation mechanism should be a three-magnon process occurring on the inhomogeneities of the sample without conservation of the magnon quasimomentum. A quantitative calculation[4] of the dependence of $\Delta\nu_k$ on the wave vector for such processes yields a weak decrease of $\Delta\nu_k$ with increasing wave vector, in qualitative agreement with our results. However, the question of which relaxation process predominates calls for a more exact quantitative analysis for each concrete substance. In particular, as indicated by Kolganov, the amplitudes of four-magnon processes may turn out to be comparable with the experimentally observed relaxation time.

3. Figure 11 shows in logarithmic coordinates the temperature dependence of $\Delta\nu_0$ extrapolated to k = 0 in the temperature range from 1.35 to 2.17°K. It is seen from the figure that in the indicated temperature interval $\Delta\nu_0 = cT^{6.4\pm0.2}$.

A weaker dependence, $\Delta\nu_0 \infty T^{4.6}$, was obtained in[12] for $CuCl_2 \cdot 2H_2O$. It is difficult to expect theoretically a simple power-law $\Delta\nu_0(T)$ dependence for processes in which the magnons collide only with one another, since the energy of the excited spin waves is of the order of k_BT in the temperature region in which our experiments were performed.

4. We turn now to the singularities of the above-threshold susceptibility. As noted above, they indicate an increase of the threshold field h_c at definite values of the wave vector or of the magnetic field. Such an increase of the threshold can be naturally attributed to interactions between the spin waves and other oscillation modes in crystal, when the spin-wave spectrum intersects the spectrum of the corresponding oscillations. We have observed experimentally two such intersection points at the fields H_1 and H_2.

It follows from the experiments described above that the position of the point of intersection corresponding to the peak H_1 does not depend on the wave vector of the spin wave. Consequently, we are dealing with some crystal oscillation mode whose energy does not depend on the wave vector. This energy, however, depends on the magnetic field and on its orientation relative to the threefold axis, as shown in Figs. 8 and 9. It is possible that this is a magnetic-impurity mode analogous to that observed in[15]. We were unable, however, to observe the absorption connected with direct excitation of this mode by an electromagnetic wave. For a final explanation of the nature of these oscillations it is necessary to investigate a number of samples with controlled impurity contents. X-ray analysis showed that our samples contained the impurities Fe (~0.35%), Ni (~0.04%) and Co (0.01–0.45%).

The elementary-excitation spectrum intersected by the spin-wave spectrum at the point corresponding to the second peak (the field H_2) is characterized by a linear dispersion law (see Fig. 7) which does not de-

pend on the magnetic field. It is natural to assume that these are phonons. Such an intersection of parametrically excited spin waves with phonons was observed in $CsMnF_3$[2].

Equating the energy and momenta of the spin waves and of the phonons and having at our disposal data on the sound velocity v, we can determine the exchange constant α from the equation

$$\alpha=\frac{v}{\pi\nu_p}\left[\left(\frac{\nu_p}{2\gamma}\right)^2-H_2(H_2+H_D)-H_\Delta^2\right]^{1/2}. \tag{4}$$

Before we use this equation, we should determine the propagation direction and the polarization of the excited phonons. For this purpose, we use the fact that when the magnetic field H is rotated in the basal plane, the position and amplitude of the peak in the field H_2 remain unchanged. Consequently, the phonons of interest to us are those whose interaction with the spin waves does not depend on the direction of the external field in the basal plane. To find such phonons let us see which terms in the magnetoelastic energy are responsible for the excitation of the phonons by the spin waves.

The equation for the elastic oscillations, describing the behavior of the phonons at small k, is[16]

$$\rho\ddot{u}_i=\partial\sigma_{ik}/\partial x_k, \tag{5}$$

where $\sigma_{ik}=\partial\mathcal{H}/\partial u_{ik}$ (summation over repeated indices is implied here and below). The Hamiltonian $\mathcal{H}$ of the system, with allowance for the magnetoelastic interaction, is written in the form

$$\mathcal{H}=\mathcal{H}_M+\mathcal{H}_{me}+\mathcal{H}_e, \tag{6}$$

where

$$\mathcal{H}_e=\alpha_{iklm}u_{ik}u_{lm},\quad \mathcal{H}_{me}=\beta_{iklm}u_{ik}l_l l_m. \tag{7}$$

in the expansion of $\mathcal{H}_{me}$ we took into account the fact that $m \ll l$. Equation (5) with allowance for (7) can be reduced to the form

$$\rho\ddot{u}_i=\alpha_{iklm}\frac{\partial^2 u_m}{\partial x_k\partial x_l}+\beta_{iklm}\frac{\partial(l_l l_m)}{\partial x_k}. \tag{8}$$

We introduce a Cartesian coordinate system in which the z axis is directed along the principal axis C_3 of the crystal and the x axis along the binary axis U_2. We assume that the magnetic field H lies in the basal plane and makes an angle φ with the axis x. In our case

$$l_0\perp H,\quad \lambda\parallel H. \tag{9}$$

Substituting in (8) the expression for l, we obtain in the right-hand side a term of the form

$$\beta_{iklm}\frac{\partial}{\partial x_k}[l_{0l}l_{0m}+(l_{0l}\lambda_m+l_{0m}\lambda_l)e^{i(\omega t-\mathbf{kr})}+\lambda_l\lambda_m e^{2i(\omega t-\mathbf{kr})}]. \tag{10}$$

We are interested in a solution in the form of a plane monochromatic elastic wave $u_i=u_{0i}e^{i(\omega t-\mathbf{kr})}$. From this, with allowance for the symmetry of the tensor β_{iklm} relative to permutation of the indices l and m, we obtain from (8) a system of three linear equations

$$(\rho\omega^2\delta_{im}-\alpha_{iklm}k_k k_l)u_m=-2i\beta_{iklm}k_k l_{0l}\lambda_m. \tag{11}$$

The concrete form of the tensors α and β is determined by the symmetry of the crystal. In crystals of the $MnCO_3$ type (symmetry group D_{3d}^6), the expressions for the elastic and magnetoelastic energies are[17]

$$\mathcal{H}_e=\tfrac{1}{2}\alpha_1u_{zz}^2+2\alpha_2(u_{xx}+u_{yy})^2+\alpha_3[(u_{xx}-u_{yy})^2+4u_{xy}^2]+2\alpha_4(u_{xx}+u_{yy})u_{zz}+4\alpha_5(u_{xz}^2+u_{yz}^2)+4\alpha_6[(u_{xx}-u_{yy})u_{yz}+2u_{xy}u_{xz}], \tag{12}$$

$$\mathcal{H}_{me}=\beta_1(u_{xx}+u_{yy})l_z^2+\beta_2u_{zz}l_z^2+\beta_3[4u_{xy}l_xl_y+(u_{xx}-u_{yy})(l_x^2-l_y^2)]+\beta_4[2u_{xz}l_xl_y+u_{yz}(l_x^2-l_y^2)]+\beta_5(u_{xx}+u_{yy})(l_x^2+l_y^2)+\beta_6u_{zz}(l_x^2+l_y^2)+\beta_7[u_{xz}l_xl_z+u_{yz}l_yl_z]+\beta_8[2u_{xy}l_xl_z+(u_{xx}-u_{yy})l_yl_z]. \tag{13}$$

After linearization with respect to λ with allowance for the second relation of (9), the expression for $\mathcal{H}_{me}$ becomes much simpler.

$$\mathcal{H}_{me}=2\lambda l_0e^{i(\omega t-\mathbf{kr})}\{\beta_3[2u_{xy}\cos2\varphi-(u_{xx}-u_{yy})\sin2\varphi]+\beta_4[u_{xz}\cos2\varphi-u_{yz}\sin2\varphi]\}. \tag{14}$$

For phonons propagating along the z axis, the system (11) reduces to

$$\begin{aligned}(C_{44}k^2-\rho\omega^2)u_x&=i\lambda l_0\beta_4k\cos2\varphi,\\(C_{44}k^2-\rho\omega^2)u_y&=-i\lambda l_0\beta_4k\sin2\varphi,\\(C_{33}k^2-\rho\omega^2)u_z&=0,\end{aligned} \tag{15}$$

where C_{33} and C_{44} are the corresponding elastic moduli. It is seen from (15) that no longitudinally polarized phonons propagating along the z axis are excited, since $u_z = 0$, and the amplitude of the transversely polarized phonons is $u_t=[u_x^2+u_y^2]^{1/2}\sim\beta_4$ and does not depend on φ. The propagation velocity of such phonons is $\nu_{tz}=\sqrt{C_{44}/\rho}$. We can show analogously that the amplitude of all the remaining phonons depends significantly on the angle φ.

Thus, it follows from our experiments and from the calculation that transversely polarized phonons propagating along the z axis are excited in the field H_2.

V. R. Gakel' measured the propagation velocity of transverse sound in $MnCO_3$ along the z axis and obtained

$$v_{tz}=3.45\cdot10^5\ [\text{cm/sec}]\pm5\%. \tag{16}$$

Using this result, we calculated from (4) the volume of the corresponding component of the constant α:

$$\alpha_z=0.79\cdot10^{-5}\ [\text{kOe-cm}]\pm7\%. \tag{17}$$

According to the data of Holden, Martel, and Svenson (private communication) on elastic scattering of neutrons in $MnCO_3$, $\alpha_z = 0.84\times10^{-5}$ kOe-cm. The spin-wave velocity obtained by us is in good agreement with this value.[2)].

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[2)]Through an error, the value of α_z in [4] is larger by a factor 2π.

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