Spin-phonon interaction and the lifetime of local vibrations in $Ni_{1-x}Be_x$

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(Submitted 24 December 1982)
Zh. Eksp. Teor. Fiz. **85**, 2032–2042 (December 1983)

The concentration and temperature dependences of the spectral distribution of local vibrations (LV) in the Ni lattice containing Be substitutional impurities have been studied by the method of inelastic slow neutron scattering. The observed nonmonotonic dependence of LV width is explained by a contribution from magnon scattering. The coupling constant for the interaction between LV and spin waves could be estimated from a quantitative comparison of the theoretical and experimental results.

PACS numbers: 63.20.Pw, 61.55.Hg, 75.30.Ds

1. INTRODUCTION

There has been much work (see, for example, Refs. 1–4) on theoretical and experimental studies of the broadening of the spectral distribution of local vibrations (LV). The main mechanisms for the broadening in nonmetallic systems for a small impurity concentration are the decay of the LV into two or more vibrations of the continuous spectrum[1,2] and the modulation mechanism, caused by scattering of phonons by LV. In metals with impurities which form LV, interaction of the local excitations with the conduction electrons is observed as well as with phonons.[5–7] Additional broadening of LV is possible in magnetic crystals associated with their interaction with the system of spin excitations, including modulation splitting of the LV spectrum due to interaction with localized spin excitations of the same impurity atom.[8]

For nonmagnetic impurities, interaction of LV with magnons of the continuous spectrum plays an important part. Such a mechanism of broadening is only effective for a sufficiently large LV amplitude of the magnetic atoms closest to the impurity, i.e., when the LV frequency lies close to the band edge of the phonon spectrum. As a consequence of the magnetic-moment conservation, the main mechanism for LV broadening due to interaction with magnons in crystals with small magnetic anisotropy is scattering of magnons by LV. The simplest and most effective scattering process is the absorption of a magnon by the local vibration with the production of a high-frequency magnon. It turns out that such a process is possible (according to the law of conservation of energy) if the LV energy ε_l is less than the maximum spin wave energy ε_m.

In the present work studies have been made of the temperature and concentration dependences of the LV width in the $Ni_{1-x}Be_x$ system to investigate the influence of various mechanisms on the LV lifetime. It is found that, unlike the case of a nonmagnetic lattice,[9] the LV width depends nonmonotonically on the Be impurity concentration and temperature. Such a behavior can be explained by the LV broadening mechanism mentioned above with the participation of two magnons. On raising the temperature (and also the concentration of nonmagnetic atoms), a softening of the magnon spectrum takes place, so that ε_m can become less than ε_l and the mechanism ceases to make a contribution to LV broadening.

2. THE EXPERIMENTS

The Ni matrix was chosen because it represents a classic ferromagnet with small magnetic anisotropy and a sufficiently high Curie temperature. In addition, the magnetic properties of Ni are well known, in particular the magnon dispersion curve has been studied experimentally.[10] Local modes produced by replacing Ni by Be were first investigated by Zemlyanov *et al.*[11] The LV line was easily observed in slow neutron inelastic scattering spectra and was characterized by a large width. However, they did not carry out a detailed study of the temperature and concentration dependences of the LV broadening.

The $Ni_{1-x}Be_x$ system has a broad region of α-solid solutions ($x \lesssim 15.3$ at. % at 1100 K), so that the magnetic characteristics of the system can be appreciably altered by changing the concentration of nonmagnetic impurities. The effect of the magnetic subsystem on the LV width can thus be studied on changing temperature and Be concentration.

$Ni_{1-x}Be_x$ alloys were studied with $x = 0.004$, 0.012, 0.025, 0.039, 0.057, and 0.092. Nickel with less than 2×10^{-4} wt. % of other impurities was used as matrix. The specimens were prepared in vacuum induction furnaces. Homogenization was carried out for 15 h at 1100 K. After quenching and tempering of the alloys produced, annealing was carried out at 650 K for up to 50 h. The method of preparing the samples was aimed at producing optimal conditions for the decay of precipitates in aged Ni–Be solid solutions.[12] Chemical analysis of specimens taken over the whole thickness of the billets showed uniformity of Be impurity distribution within 5%. Metallographic, x-ray and neutron structural analyses, and also the results of slow neutron small angle scattering showed that the specimens investigated were disordered α-solutions of Be in Ni. Measurement was made of the intensity of optic modes ($\varepsilon = 65$ meV) in the spectrum of inelastic neutron scattering for a specially prepared two-phase $Ni_{82}Be_{18}$ specimen, containing 9 at. % Be in the intermetallic β-phase and for the α-solutions studied, in order to estimate possible small amounts of intermetallic inclusions, which are difficult to observe by traditional methods. By this means it could be established that the possible intermetallic content in the specimens studied ($x = 0.004$–0.092) is negligible. For the most enriched beryllium specimen there is less than 0.2 at. % Be in the β-phase.

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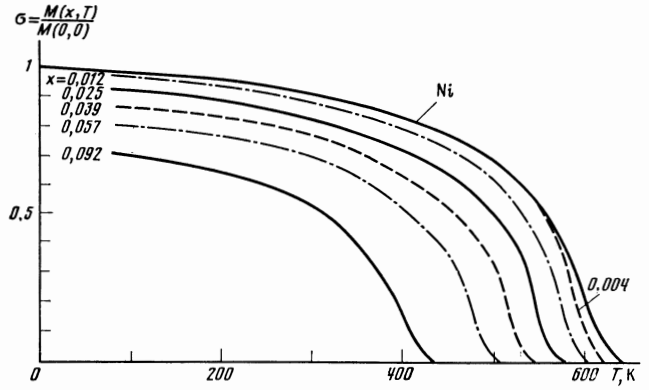

FIG. 1. Temperature dependence of the magnetization of the $Ni_{1-x}Be_x$ alloys studied.

The results of measurement of the magnetization of the specimens are shown in Fig. 1. It follows from the results, in particular, that the Curie temperature decreases according to the law $T_c(x) = T_c(0)(1-bx)$, where $b = 3.51 \pm 0.06$, while the magnetization at $T \to 0$ varies as $M(x,0) = M(0,0)(1-3.2x)$. Measurements of LV spectra were carried out in the temperature range 170–520 K (below the temperature T_c of the phase transition to the nonmagnetic state) to study the effect of the magnetic subsystem on the LV lifetime. It can be seen in Fig. 1 that in just this range, a noticeable change takes place in the magnetization of the specimens with various Be concentrations.

Measurements of LV spectra were carried out by means of slow neutron inelastic scattering. The doubly differential neutron scattering cross-sections $d^2\sigma/d\varepsilon d\Omega$ were measured by a time-of-flight spectrometer.[13] The spectral distribution of Be LV in the Ni lattice could be extracted from measurements of scattering cross-sections, with the same geometry and neutron transmission, on pure Ni and Ni–Be (Fig. 2). Measurements were carried out on polycrystalline specimens and analysis of the neutron cross-sections for scattering by LV was carried out in the noncoherent approximation.

The values of natural LV spectral widths, determined from the spectral distributions obtained, are shown in Fig. 3 together with the experimental error. The natural LV widths were found by fitting, to the experimental spectra, convolutions of the Lorentzian function that describes the LV line shape sought with the experimentally measured[14] resolution function of the spectrometer. The error in the values shown in Fig. 3 was determined by taking into account the errors in determining the energy width of the spectrometer and the LV spectral distributions, averaged overall measured angles Ω_i.

The concentration and temperature dependences of LV widths, studied earlier,[9,15] for solutions of Be substituted in copper and vanadium lattices, are also shown in Fig. 3 for comparison. As can be seen from the figure, the linear concentration and monotonic temperature dependences of the LV width, characteristic of a nonmagnetic matrix, are only observed for Ni–Be in the low impurity concentration range and at low temperatures. As x and T increase, a noticeable

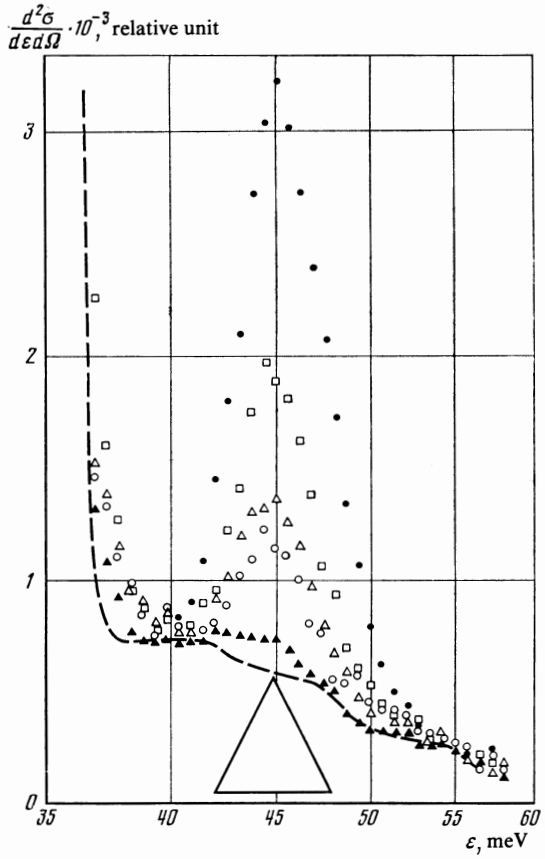

FIG. 2. Scattering cross-sections for slow neutrons by local vibrations in $Ni_{1-x}Be_x$ at room temperature and for one of the scattering angles Ω, for $x =$ ▲—0.004; ○—0.025; △—0.039; □—0.057 and ●—0.092; the dashed curve is for scattering by pure Ni. The cross sections have been normalized to unit neutron flux and the same number of matrix atoms in the scattering systems.

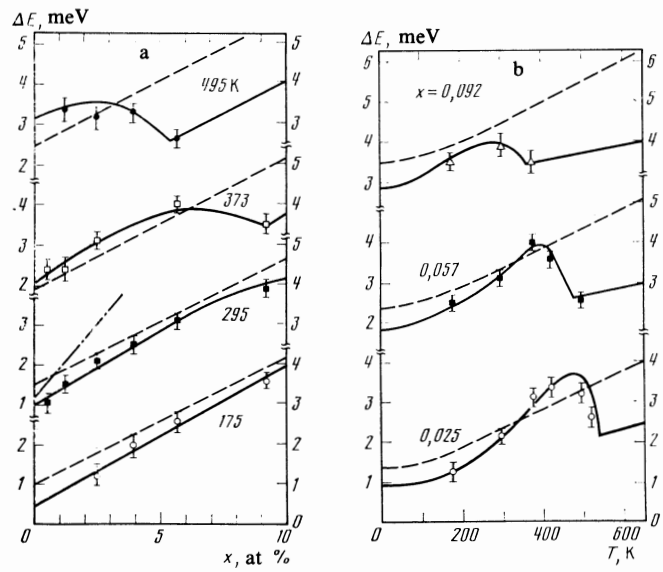

FIG. 3. Concentration (a) and temperature (b) dependences of the width of LV spectra. Points—experiment, full lines—theory; the dashed lines show the concentration and temperature dependences for LV width in Cu–Be[7] and the dash-dot line is for V–Be[15].

departure from linearity occurs. The nonmonotonic behavior of LV width takes place, naturally, in the T and x region where the magnetic properties of the system change appreciably (see Fig. 1). In particular, the position of the feature in the temperature dependence (Fig. 3b) correlates with the temperature region where the magnetization of the alloys studied decreases sharply. This indicates the existence of an additional mechanism for LV broadening in magnetic systems. The interaction of LV with magnons may be such a mechanism.

3. THEORY

As has already been pointed out, the line width of any nonmagnetic excitation in an isotropic magnetic system, due to interactions with the magnetic subsystem, can only be associated with processes for magnon scattering by these excitations. The most intense of these processes, if it is allowed by the law of conservation of energy, is scattering with the participation of the minimum number (i.e. two) of magnons. The LV broadening produced by such a process has the form

$$\Gamma_M \equiv \Gamma_M(\varepsilon_l) = \pi \sum_{\varkappa_1 \varkappa_2} |\gamma_{\varkappa_1 \varkappa_2}|^2 (n_{\varkappa_1} - n_{\varkappa_2}) \delta(\varepsilon_l - \varepsilon_{\varkappa_2} + \varepsilon_{\varkappa_1}),$$
$$n_\varkappa \equiv n(\varepsilon_\varkappa) = [\exp(\varepsilon_\varkappa / kT) - 1]^{-1}, \qquad (1)$$

where $\gamma_{\varkappa_1 \varkappa_2}$ is the magnon-phonon interaction matrix element and $\varepsilon_\varkappa$ is the magnon energy. If the value of $\gamma_{\varkappa_1 \varkappa_2}$ can be represented in the form of a function of $\varepsilon_{\varkappa_1}$ and $\varepsilon_{\varkappa_2}$, then Eq. (1) takes the form

$$\Gamma_M(\varepsilon_l) = \int dE_1 \int dE_2 \tilde{\gamma}^2(E_1, E_2) g(E_1) g(E_2) \times [n(E_1) - n(E_2)] \delta(\varepsilon_l - E_2 + E_1). \qquad (2)$$

Here $g(E)$ is the magnon continuous spectrum density of states.

We will, for simplicity, discuss Eq. (2) within the framework of the Heisenberg model. We will then write the Hamiltonian of a ferromagnet with a nonmagnetic impurity atom in the following form:

$$H = H_s + H_p; \qquad H_s = -\frac{1}{2} \sum_{\mathbf{nn}'} \mathbf{S_n} I_{\mathbf{nn}'} \mathbf{S_{n}}';$$
$$H_p = \sum_k \varepsilon_k b_k^+ b_k + H_a. \qquad (3)$$

Here H_s is the spin Hamiltonian of the ferromagnet, H_p is the elastic wave (phonon) Hamiltonian, H_a is the anharmonicity Hamiltonian; $I_{nn'}$ is the exchange energy of a pair of spins at sites $\mathbf{n}$ and $\mathbf{n}'$ ($n, n' \neq 0$, since the nonmagnetic impurity atom is placed at the site with index 0); $\mathbf{S}_n$ is the atomic spin matrix operator; ε_k is the energy of a phonon with wavenumber k (the oscillations in a crystal containing impurities are not described by plane waves), where k includes both LV ($k = l$) and oscillations of the continuous spectrum.

We shall separate in the Hamiltonian H_s the interaction of oscillations with the spin system, expanding the exchange integral $I_{nn'}$ in powers of small displacements of the atoms from their equilibrium positions

$$\mathbf{u_n} = \sum_k \alpha_{k\mathbf{n}} b_k + \text{H.c.}$$

and limit ourselves to the linear term in the expansion. It can be shown by using, for example, the equal-time Green function, that in second order perturbation theory in the spin-phonon interaction, the broadening of the LV spectral distribution has the form

$$\Gamma_M = \{2[1 + n(\varepsilon_l)]\}^{-1} \sum_{\mathbf{nn_1 n_2 n_3}} B_{l\mathbf{nn_1}} B^*_{l\mathbf{n_2 n_3}} \langle \mathbf{S_n S_{n_1}}; \mathbf{S_{n_2} S_{n_3}} \rangle_{\omega = \varepsilon_l}, \qquad (4)$$

where

$$B_{l\mathbf{nn_1}} = (\alpha_{l\mathbf{n}} - \alpha_{l\mathbf{n_1}}) \nabla I_{\mathbf{nn_1}},$$
$$\langle A; B \rangle_\omega = \frac{1}{2\pi} \int_{-\infty}^{\infty} \exp(i\omega t) \langle A(t) B(0) \rangle dt,$$

the angle brackets signify quantum-statistical averaging.

At temperatures not too near to T_c, when $(T_c - T)/T_c \gg Z^{-2}$, where Z is the coordination number, a self-consistent field approximation of the type considered by Vaks *et al.*[16] can be used for the calculation of the correlation function in Eq. (4). In this approximation (if processes involving four and more magnons are neglected), Eq. (4) can be reduced to the form of Eq. (1) with

$$\varepsilon_\varkappa = \sigma E_\varkappa, \qquad E_\varkappa = S[I(0) - I(\varkappa)],$$
$$\sigma = \bar{S}/S, \qquad \bar{S} = \bar{S}(T, x) = \frac{1}{N} \sum_{\mathbf{n}} \langle \mathbf{S_n}^2 \rangle,$$
$$S = S(0,0); \qquad I(\varkappa) = \sum_{\mathbf{n}'} \exp(i\varkappa \mathbf{n}') I_{\mathbf{nn}'}; \qquad (5)$$
$$\gamma_{\varkappa_1 \varkappa_2} = \frac{1}{2} \bar{S} \sum_{\mathbf{nn}'} B_{l\mathbf{nn}'} (A_{\varkappa_1 \varkappa_2 \mathbf{nn}} + A_{\varkappa_1 \varkappa_2 \mathbf{n}'\mathbf{n}'} - A_{\varkappa_1 \varkappa_2 \mathbf{nn}'} - A_{\varkappa_1 \varkappa_2 \mathbf{n}'\mathbf{n}}),$$
$$A_{\varkappa_1 \varkappa_2 \mathbf{nn}'} = \frac{1}{N} \exp[i(\varkappa_1 \mathbf{n} - \varkappa_2 \mathbf{n}')].$$

The quantity $\bar{S}$ here, which determines the magnetization of the system, depends in general both on T and x (see Fig. 1), N is the number of lattice sites, while the expression for $A_{\varkappa_1 \varkappa_2 nn'}$ is written with neglect of the change in spin-wave amplitude in the neighborhood of the impurity.

Corresponding to the experimental results for the Ni–Be system, we consider the case when at $T = 0$, $\varepsilon_l < \varepsilon_m$ and the corresponding process is allowed over a fairly wide temperature range. At low temperatures when $kT \ll \varepsilon_m - \varepsilon_l$, the main contribution to the value of Γ_m (Eq. 1) comes from the region where the energy of the absorbed magnon $\varepsilon_{\varkappa_1} \sim kT$. Expanding Eq. (5) in terms of small $\varkappa_1$, we obtain

$$\varepsilon_{\varkappa_1} = \alpha_1 (\varkappa_1 d)^2 \varepsilon_m \ (\varkappa_1 d \ll 1), \ \varepsilon_m = E_m \sigma;$$
$$\gamma_{\varkappa_1 \varkappa_2} = N^{-1} \bar{S} \varkappa_1 \cdot \frac{1}{2} \sum_{\mathbf{nn}'} B_{l\mathbf{nn}'} |\mathbf{n} - \mathbf{n}'| \cos \beta_1 (e^{i\varkappa_2 \mathbf{n}} - e^{i\varkappa_2 \mathbf{n}'}), \qquad (6)$$

where d is the period of the cubic lattice, β_1 is the angle between the vectors $\varkappa_1$ and $\mathbf{n} - \mathbf{n}'$. For the case of Ni with an fcc lattice, which interests us, in the nearest-neighbor approximation $\alpha_1 = 1/16$, and the limiting spin-wave frequen-

cy at $T=0$ in an ideal crystal $E_m = 16\,IS$,[17] where I is the exchange integral for nearest neighbors of the matrix.

Substituting Eq. (6) in Eq. (1) we then obtain

$$\Gamma_M = A(kT/E_m)^{5/2}(kT \ll E_m - \varepsilon_l),$$
$$A = \frac{5}{64}\frac{1}{\sqrt{\pi}}\xi\left(\frac{5}{2}\right)S^2\frac{v}{d^5}\frac{1}{\alpha_1^{5/2}}\frac{1}{N}\sum_{\varkappa_2} B_{\varkappa_2}B_{\varkappa_2}^{*}\delta(\varepsilon_{\varkappa_2}-\varepsilon_l), \tag{7}$$
$$B_{\varkappa_2} = \sum_{nn'} B_{lnn'}(\mathbf{n}-\mathbf{n}')[\exp(i\varkappa_2\mathbf{n}) - \exp(i\varkappa_2\mathbf{n}')].$$

The temperature dependence of the broadening obtained $\propto (kT)^{5/2}$ at low temperatures is only a consequence of the spin-wave quadratic dispersion law and does not depend on the choice of model approximations. At very low temperatures the contribution from this mechanism can be neglected compared with broadening due to anharmonicity.[1]

As the temperature is raised, the broadening determined by Eq. (1) grows rapidly and for $kT \gtrsim \varepsilon_m - \varepsilon_l$ can be the main LV broadening mechanism. For analyzing the temperature and concentration dependences of this mechanism at sufficiently high temperature, we approximate the value of $\gamma_{\varkappa_1\varkappa_2}$ over the whole spin-wave region of the spectrum by its value corresponding to small $\varkappa_1$ and $\varkappa_2$. Expanding Eq. (6) (valid for small $\varkappa_1$) in terms of $\varkappa_2$, we obtain

$$\gamma_{\varkappa_1\varkappa_2} = N^{-1}\gamma_0\sigma(\varepsilon_{\varkappa_1}/\varepsilon_m)^{1/2}(\varepsilon_{\varkappa_2}/\varepsilon_m)^{1/2},$$
$$\gamma_0 = \frac{1}{2}\sum_{nn'}\frac{S}{\alpha_1 d^2}B_{lnn'}|\mathbf{n}-\mathbf{n}'|^2\cos\beta_1\cos\beta_2, \tag{8}$$

where β_2 is the angle between the vectors $\varkappa_2$ and $\mathbf{n}-\mathbf{n}'$. We note that other physically acceptable approximations for $\gamma_{\varkappa_1\varkappa_2}$ as shown by numerical analysis, do not change qualitatively the results obtained below. Substituting Eq. (8) into Eq. (1), we obtain the expression for Γ_M in the form of Eq. (2), in which

$$\tilde{\gamma}(E_1, E_2) = \pi^{1/2}\gamma_0\sigma(E_1/\varepsilon_m)^{1/2}(E_2/\varepsilon_m)^{1/2}. \tag{9}$$

It is convenient, for what follows, to write the expressions for the density of states in the following form:

$$g(E) = \varepsilon_m^{-1}\tilde{g}(E/\varepsilon_m), \quad \tilde{g}(y) = N^{-1}\sum_{\varkappa}\delta(y-\varepsilon_\varkappa/\varepsilon_m). \tag{10}$$

In the self-consistent field approximation, the dimensionless magnon density of states of the crystal, $\tilde{g}(y)$, is independent of temperature and concentration, since in this approximation $\varepsilon_\varkappa/\varepsilon_m = E_\varkappa/E_m$, and the sum over $\varkappa$ in Eq. (10) is carried out over the whole Brillouin zone. Substituting Eqs. (9) and (10) in (2), we obtain

$$\Gamma_M(T, x) = \gamma\varepsilon_l f(T, x),$$
$$f(T,x) = \sigma\int_0^{1-a/\sigma} dy\,y\left(y+\frac{a}{\sigma}\right)\tilde{g}(y)\tilde{g}\left(y+\frac{a}{\sigma}\right)\left[n(E_m\sigma y) - n\left(E_m\sigma\left(y+\frac{a}{\sigma}\right)\right)\right], \quad \gamma = \pi\gamma_0^2/\varepsilon_l E_m, \quad a = \varepsilon_l/E_m. \tag{11}$$

Strictly speaking, Eq. (11) is valid only for small impurity concentration. It can be shown that Eq. (11) also stays valid for somewhat larger x if σ is taken in the form $\sigma(T,x) = M(T,x)/M_0$, where $M(T,x)$ is the experimentally determined magnetization of the crystal for given T and x, while $M_0 \equiv M(0,0)$. The main concentration dependence is then taken into account and the dependence of the form $1-x$ is omitted, describing in particular the change in the integral spin-wave density of states on introducing nonmagnetic impurities.

The LV broadening produced by spin-phonon interaction in a ferromagnet with nonmagnetic impurities can be determined from Eq. (11) over a sufficiently wide range of variation of temperature and impurity concentration required for comparison with the experimental results.

It can be seen from the expression that at low temperatures $kT < \varepsilon_m - \varepsilon_l$, the broadening $\Gamma_M \propto T^{5/2}$, corresponding to the result of Eq. (7). However, at high temperatures, the growth in Γ_M with increasing temperature is replaced by a fall. This is associated with the fact that at sufficiently high temperatures the magnetization of the matrix σ can decrease appreciably, the maximum frequency in the spin-wave spectrum ε_m, as can be seen from Eqs. (5) and (6), also falls and the region of allowed values $\varepsilon_m - \varepsilon_l$ of the spin wave energy [and correspondingly the region for integrating over y in Eq. (11)] is considerably contracted. If we do not consider temperatures too close to T_c and we neglect the contribution of relaxation magnetic excitations to the LV broadening, the broadening produced by two-magnon scattering by LV vanishes at a temperature for which $\varepsilon_m = \varepsilon_l$. As the analysis shows, such a behavior of Γ_M is also preserved in a more general case (without using the molecular-field approximation) and is a consequence of softening of the spectrum of the magnetic excitations of the system with increasing T.

We shall give, as an example, the appropriate $\Gamma_M(T,x)$ dependences for various x for the Ni–Be system studied, for which $E_m \approx 140$ meV, according to experimental data,[10] and $\varepsilon_l = 44$ meV. We choose the density of states $\tilde{g}(E/\varepsilon_m)$ in the

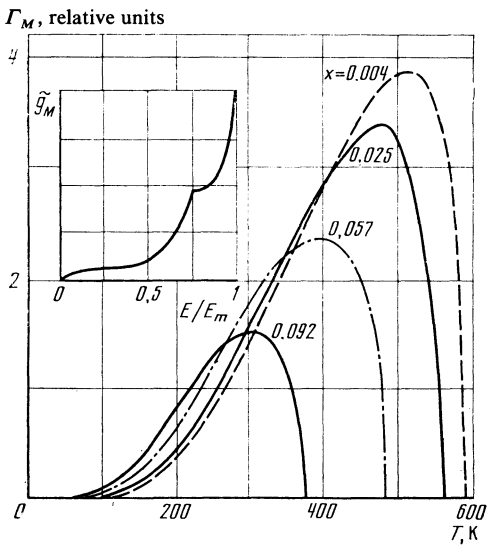

FIG. 4. Temperature dependence of the contribution Γ_M to the LV broadening for $Ni_{1-x}Be_x$. The values of x are shown for each curve. Inset: model[17] Ni magnon spectrum.

form corresponding to an fcc lattice with nearest neighbor interaction.[17] Using the experimental values of $\sigma(T,x)$ (Fig. 1) we obtain the $\Gamma_M(T,x)$ dependences shown in relative units in Fig. 4. It can be seen that with increasing T the dependence of Γ_M on T goes through a maximum and that at $\varepsilon = \varepsilon_l$ it returns to zero with the corresponding temperature less than the Curie temperature, so that the system maintains its long-range magnetic order. On the other hand, Γ_M decreases on increasing the concentration of nonmagnetic impurities. The appearance of the curves shown remains practically unchanged if some other approximation is used for $\tilde{\gamma}(E_1,E_2)$ and $\tilde{g}(E,\varepsilon_m)$. In the next section these results are used to explain the experimental data.

4. DISCUSSION OF THE RESULTS

Apart from the broadening Γ_M produced by the magnetic subsystem, there are contributions to the experimentally observed LV width in the Ni–Be system associated with a number of other mechanisms which are also characteristic of nonmagnetic systems. The main such mechanism for small impurity concentrations is the anharmonic coupling between LV and phonons which leads to processes of decay of LV into two or more phonons[1,2] and also to modulation scattering of phonons by LV.[3] If the energy conservation law allows the decay of LV into two phonons of the continuous spectrum, then such a process is dominant at not too high temperatures. Its temperature dependence can be described by the following interpolation formula:[3]

$$\Gamma_a = \frac{\alpha}{4}\varepsilon_l\left[1+2n\left(\frac{\varepsilon_l}{2}\right)\right], \tag{12}$$

where α is a dimensionless parameter characterizing the anharmonic LV-phonon interaction and $n(\varepsilon)$ is the level population of lattice excitations.

In metals, LV broadening can also be associated with their interaction with conduction electrons.[5] However, as the study of several similar (nonmagnetic) systems with Be impurity shows,[7] this interaction can be neglected in the present case.

Broadening of LV in nonideal crystals can also be produced by interaction between the impurities. Dynamic interaction between LV resulting from exchange of virtual phonons depends exponentially on the distance between impurities, so that the corresponding broadening depends exponentially on concentration.[18] In the system considered, the LV frequency lies far from the band edge of the continuous spectrum, and the radius of the localized excitation is consequently $\sim v^{1/3}$. For not too high concentrations, therefore, the overlap of the LV wave functions of different centers is small and the corresponding LV broadening produced by dynamic interaction can also be neglected.

On the other hand, the existence of impurities leads to the appearance of long-range elastic deformations in the crystal. These deformations randomly shift and split the LV line, leading to its inhomogeneous broadening. If it is assumed that these deformations are produced by the same Be impurity atoms as lead to the appearance of LV, then the concentration broadening in the alloys studied can be written in the form[19]

$$\Gamma_x = \beta\varepsilon_l x \quad (x \leqslant 0.1), \tag{13}$$

where β is a dimensionless coefficient which depends on the change in volume of the crystal produced by a defect and on the parameters of the anharmonic impurity-matrix interaction.

The independence of α of x and of β of T (i.e. the additivity of Γ_α and Γ_x) in the temperature and concentration range studied is confirmed experimentally.[7] We will therefore assume that the total LV width Γ_l in the $Ni_{1-x}Be_x$ system can be described by the following expression:

$$\Gamma_l = \Gamma_a + \Gamma_x + \Gamma_M. \tag{14}$$

We will thus use three dimensionless adjustable parameters α, β, γ to describe the experimental values of $\Gamma_l(x,T)$. Their values for Ni–Be obtained by the least squares method are shown in Table I. Curves corresponding to these parameters are shown by dashed lines in Fig. 3. We also show in the table values of $\xi = 1 - m_{Be}/M$ (m_{Be} and M are the atomic masses of the impurity and of the appropriate matrix), the positions of the energy of an isolated (i.e. $x \to 0$) Be impurity LV, ε_l, and the calculated LV energies ε_l^i obtained in the isotopic model by using the experimentally determined phonon spectrum of the matrix.

As can be seen from the table, the values found for the parameters α and β for Ni–Be are of the same order as the corresponding quantities for Cu–Be and V–Be. A direct dependence can be traced of the parameter β, which determines the concentration LV broadening, on the quantity $\tau = (\varepsilon_l/\varepsilon_l^i)^2$, which characterizes the difference in the interaction of an isolated impurity with the matrix atoms surrounding it, compared with the interatomic interaction in the matrix. For LV in Ni and Cu, in which the effective interaction between the Be impurity atoms and the matrix is weakened appreciably, the parameters β are thus similar and considerably different from the corresponding value for V–Be, where there is practically no difference between the impurity–matrix and intramatrix force constants. It follows from this that the concentration LV broadening in comparable systems is produced by deformation fields arising from the impurity centers, as has been indicated earlier.[13]

Because the LV amplitude at matrix atoms near the impurity center is not zero, there is also an influence of the anharmonicity of the matrix lattice on the magnitude of the temperature broadening [see Eq. (12)] as well as the impurity atom-matrix interaction. In fact, for roughly equal values of τ, characterizing the relative interaction of Be impurity atoms in Ni and Cu matrices, the value of the parameter α for the Ni–Be system is less than for Cu–Be. This reflects the actual relation between the coefficients of the anharmonicity of the lattice vibrations which are manifest, in particular, in the coefficients of thermal expansion of the crystal lattices of these alloys.

The interaction of LV with spin waves in $Ni_{1-xc}Be_x$ makes an appreciable contribution to the total width Γ_l. For example, for $x = 0.05$ Γ_M at room temperature is three times greater than the decay due to anharmonicity Γ_α and

TABLE I. Values of the parameters characterizing the attenuation of local vibrations of Be in various metal matrices.

Matrix	ξ	ε_l, meV ($T \approx 300$ K)	ε_l^i, meV	α	β	γ	References
Cu	0.858	40	56	0.027±0.004	0.81±0.06 *	–	[7], [9], [19]
V	0.823	54	56.5	0.023±0.004	1.9±0.3	–	[15]
Ni	0.846	44±1	63	0.016±0.007	0.70±0.12	0.066±0.016	–

*The improvement in the accuracy, compared with earlier values,[9] results from taking account of additional measurements.[7,19]

comparable with the concentration broadening Γ_x. Such values of Γ_M imply a considerable LV amplitude at magnetic atoms nearest to the impurities and a fairly large value of magnon-phonon interaction.

It is interesting to note that additional information on the spin-wave spectrum of the magnetic sublattice can be obtained from analysis of the concentration and temperature dependences of the LV width. Using the value of $a = \varepsilon_l / E_m$ as yet another adjustable parameter, we obtained values of the edge of the model spin excitation spectrum shown in Fig. 4, $E_m = 145 \pm 18$ meV, which agrees with experimental results of measuring the dispersion law for spin waves in Ni.[10]

The analysis of the results of the present work thus show that interaction with spin excitations makes an appreciable contribution to LV broadening in magnetic materials. This contribution has a characteristic temperature dependence with a maximum. Such a behavior of Γ_M is independent of the approximations of the models used above and is described by the approximate Eq. (11).

The authors thank Yu. G. Chukalkin for help with the measurements of the magnetization of the specimens studied.

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Translated by R. Berman