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Dynamic damping of dislocations by magnetizations in ferromagnets

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The dynamic damping of dislocations by magnetizations in ferromagnets is considered. The temperature dependence of the damping force is investigated in the vicinity of the Curie point within the framework of the molecular-field approximation. The relative efficiencies of the phonon, electron, and magnetic contributions to the damping force are compared. It is shown that the magnetic contribution is decisive only in ferromagnets with low Curie points.

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It is by now reliably established that the plastic properties of metals are connected in one way or another with the motions of dislocations that are either pinned by static obstacles (impurities, complexes or impurities, other dislocations, long-range fields or elastic stresses, and others), or are dynamically damped by phonons, electrons, and other elementary excitations. Dynamic damping is decisive at high dislocation velocities, when the motion is above the barrier. It must be taken into account also in the case of low velocities when dislocations glide between obstacles. The most effective among the dynamic mechanisms is a rule damping by phonons. Contributions of other elementary excitations manifest themselves only in some distinct situations: thus, dislocation damping by electrons is significant at sufficiently low temperatures, and damping by magnons takes place in ferromagnets below the Curie point.

Dynamic damping of dislocations by phonons and electrons has been studied quite thoroughly, and is the subject of exhaustive reviews.[1,2] There is still no corresponding description of dislocation damping by magnons in ferromagnets. Among the latest papers connected with this problem mention should be made of the article by Bar'yakhtar and Druinskiĭ, who investigated dislocation damping by magnons and obtained the temperature dependence of the magnon damping

$$F_s = B_s^{(2)} (T/\Theta_c)^{5/2} V, \tag{1}$$

where V is the dislocation, $B_s^{(2)}$ is a coefficient independent of the temperature T, and Θ_c is the Curie temperature. Bar'yakhtar and Druinskiĭ[3] used the spin-wave approximation, which describes ferromagnets at relatively low temperatures. Equation (1) therefore describes correctly the temperature dependence of the magnetic-damping force at low values of T, but is not

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valid at T close to Θ_c. This is clear from the fact that at $T=\Theta_c$ a phase transition into a nonmagnetic state occurs, so that F_s should vanish.

The present paper is devoted to a study of dislocation damping by magnons in ferromagnets at temperature close to Θ_c. It is known that in this temperature region a ferromagnet is described by the molecular-field approximation, in the framework of which the calculation will be made here. In addition, we use in the calculation the following models and approximations: we consider the damping of a linear edge dislocation moving at constant velocity in a ferromagnetic crystal with a primitive cubic lattice (PCL), and describe the ferromagnet by using the Heisenberg model and the nearest-neighbor approximation.

We use a system of units with $\hbar = k = 1$.

1. THE HAMILTONIAN

We write down in general form the ferromagnet Hamiltonian in the Heisenberg model:

$$\mathcal{H} = -g\mu_0 H \sum_l S_l^{(3)} - \frac{1}{2}\sum_{l \neq l'} J(\mathbf{R}_{ll'})\mathbf{S}_l\mathbf{S}_{l'}, \tag{2}$$

where g is the Lande factor, μ_0 is the Bohr magneton, H is the constant external field, $\mathbf{S}_l$ is the spin operator at the site l ($S_l^{(3)}$ is its x_3-projection), $J(\mathbf{R}_{ll'})$ is the exchange integral and decreases rapidly (exponentially) with distance,[1)] and $\mathbf{R}_{ll'}$ is the distance between the instantaneous positions of the atoms:

$$\mathbf{R}_{ll'} = \mathbf{R}^0_{ll'} + \mathbf{u}_{ll'}. \tag{3}$$

Here $R^0_{ll'}$ is the distance between the lattice sites l and l', and $\mathbf{u}_{ll'} = \mathbf{u}_{l'} - \mathbf{u}_l$ is the relative displacement of the atoms.

To describe the interaction of a dislocation with a magnon, we assume that the displacements $\mathbf{u}_{ll'}$ are due to the dislocation (in our case, one edge dislocation). In the nearest-neighbor approximation $\mathbf{u}_{ll'} \ll a$, where a is the lattice constant. We can therefore expand the exchange integral in terms of small relative displacements and confine ourselves to the terms linear in $\mathbf{u}_{ll'}$:

$$J(\mathbf{R}_{ll'}) = J(\mathbf{R}^0_{ll'}) + \sum_\mu \nabla^{(\mu)}_{ll'} J(\mathbf{R}_{ll'}) u^{(\mu)}_{ll'} + \ldots \tag{4}$$

Here

$$\nabla^{(\mu)}_{ll'} J(\mathbf{R}_{ll'}) = \frac{\partial J(R_{ll'})}{\partial R^{(\mu)}_{ll'}}\Big|_0$$

and the summation is over the Cartesian coordinates ($\mu = 1, 2, 3$).

When (4) is taken into account, we can write the Hamiltonian in the nearest-neighbor approximation for a PCL in the form

$$\mathcal{H} = -g\mu_0 H \sum_l S_l^{(3)} - \frac{J}{2}\sum_{l\Delta} \mathbf{S}_l\mathbf{S}_{l+\Delta} - \frac{1}{2}\sum_l \sum_{\mu\beta} \sum_\varkappa \nabla^{(\mu)}_{l,l+\varkappa\Delta_\beta} J(\mathbf{R}^0_{l,l+\varkappa\Delta_\beta}) u^{(\mu)}_{l,l+\varkappa\Delta_\beta} \mathbf{S}_l\mathbf{S}_{l+\varkappa\Delta_\beta}. \tag{5}$$

The summation over Δ means summation over the neighbors nearest to the site l; the index $\varkappa$ can take on two values, plus or minus, $J \equiv J(\Delta) = J(\mathbf{R}^0_{l,l+\Delta})$. In the PCL, Δ has six components:

$$\pm\Delta_1 = (\pm\Delta_1, 0, 0); \quad \pm\Delta_2 = (0, \pm\Delta_2, 0); \quad \pm\Delta_3 = (0, 0, \pm\Delta_3).$$

We shall hereafter denote the first two terms in the right-hand side of (5) by $\mathcal{H}_s$, and the next ones by $\mathcal{H}_{sd}$. It is convenient to change over in $\mathcal{H}_{sd}$ from relative displacements to strains. We do this with the relation

$$u^{(\mu)}_{l,l+\varkappa\Delta_\beta} = \frac{u^{(\mu)}_{l+\varkappa\Delta_\beta} - u^{(\mu)}_l}{a} a \approx \varkappa a \frac{\partial u^{(\mu)}_l}{\partial x_\beta}. \tag{6}$$

Taking (6) into account we rewrite $\mathcal{H}_{sd}$ in the form

$$\mathcal{H}_{sd} = -\frac{a}{2}\sum_l \sum_{\mu\beta} \sum_\varkappa \varkappa \nabla^{(\mu)}_{l,l+\varkappa\Delta_\beta} J(\mathbf{R}_{l,l+\varkappa\Delta_\beta}) \frac{\partial u^{(\mu)}_l}{\partial x_\beta} \mathbf{S}_l\mathbf{S}_{l+\varkappa\Delta_\beta}. \tag{7}$$

We consider now the concrete case of the damping of an edge dislocation directed along the x_3 axis and having a Burgers vector $\mathbf{b} = (b, 0, 0)$. For this dislocation, there are two nonzero components, $u^{(1)}_l$ and $u^{(2)}_l$, i. The index μ of (7) takes on only two values (1 and 2). The exchange integral $J(\mathbf{R})$ depends in fact on the modulus $|\mathbf{R}|$, the main dependence on $|\mathbf{R}|$ being of the form (see, e.g., Ref. 4)

$$J(\mathbf{R}) \sim e^{-2|\mathbf{R}|/\alpha}, \tag{8}$$

where $\alpha = \hbar^2/me^2 \approx 5\times 10^{-9}$ cm. It is easy to show that within the framework of the approximations assumed by us for an edge dislocation the only terms remaining in the quantities

$$\nabla^{(\mu)}_{l,l+\varkappa\Delta_\beta} J(\mathbf{R}_{l,l+\varkappa\Delta_\beta})$$

are those with equal indices μ and β, and

$$\nabla^{(\mu)}_{l,l+\varkappa\Delta_\mu} J(\mathbf{R}_{l,l+\varkappa\Delta_\mu}) \approx -2\varkappa J/\alpha. \tag{9}$$

Consequently, $\mathcal{H}_{sd}$ for an edge dislocation can be represented in the form

$$\mathcal{H}_{sd} = \frac{aJ}{\alpha}\sum_l \sum_{\mu\varkappa} \varepsilon_{\mu\mu}(l)\mathbf{S}_l\mathbf{S}_{l+\varkappa\Delta_\mu}, \tag{10}$$

where $\varepsilon_{\mu\mu}(l)$ is a component of the strain tensor.

We expand the strain tensor in a Fourier series. We recognize here that the dislocation, and with it the strain field, moves with constant velocity V:

$$\varepsilon_{\mu\mu}(\mathbf{l} - \mathbf{V}t) = \frac{1}{N}\sum_{\mathbf{q}} \varepsilon_{\mu\mu}(\mathbf{q}) e^{i\mathbf{q}\mathbf{l}} \exp(-i\Omega_{\mathbf{q}} t), \tag{11}$$

where $\Omega_{\mathbf{q}} = \mathbf{q}\cdot\mathbf{v}$. Taking (11) into account we write down the final expression for the Hamiltonian of the interaction of a dislocation with a magnon:

$$\mathcal{H}_{sd}(t) = \frac{1}{N}\sum_{\mathbf{q}\mathbf{l}} \sum_{\mu\varkappa} \Phi_{\mu\mu}(\mathbf{l}, \mathbf{q}) \exp(-i\Omega_{\mathbf{q}} t)\mathbf{S}_l\mathbf{S}_{l+\varkappa\Delta_\mu}, \tag{12}$$

where

$$\Phi_{\mu\mu}(\mathbf{l}, \mathbf{q}) = aJ\varepsilon_{\mu\mu}(\mathbf{q}) e^{i\mathbf{q}\mathbf{l}}/\alpha. \tag{13}$$

2. DISLOCATION-ENERGY ABSORPTION PER UNIT TIME

Following Al'shitz and Mal'shukov,[5] we calculate the magnetic component of the absorption of energy by a moving dislocation in a ferromagnet, using the linear-response theory[6]:

$$D=-\mathrm{Sp}\left[\Delta\rho(t)\frac{\partial\mathcal{H}_{sd}(t)}{\partial t}\right], \tag{14}$$

where $\Delta\rho(t)=\rho(t)-\rho_0$, and ρ_0 denotes the equilibrium density matrix of a crystal without dislocations.

From the equation for the density matrix

$$\frac{\partial\rho}{\partial t}+i[\mathcal{H},\rho]=0 \tag{15}$$

we obtain in the approximation linear in the perturbation $\mathcal{H}_{sd}$

$$\Delta\rho(t)=-i\int_{-\infty}^{t}dt_1\exp\{i\mathcal{H}_s(t_1-t)\}[\mathcal{H}_{sd}(t_1),\rho_0]\exp\{-i\mathcal{H}_s(t_1-t)\}. \tag{16}$$

From (14), taking (16) and (12) into account, it follows that

$$=\frac{i}{N^2}\sum_{\substack{\mathbf{q}\lambda\mu\\ \mathbf{q}_1\lambda_1\mu_1}}\Phi_{\mu\mu}(\mathbf{l},\mathbf{q})\Phi_{\mu_1\mu_1}(\mathbf{l}_1,\mathbf{q}_1)\Omega_{\mathbf{q}_1}\exp\{-i(\Omega_{\mathbf{q}}+\Omega_{\mathbf{q}_1})t\}G^{r}_{\mathbf{l}_1,\mathbf{l}_1+\varkappa_1\Delta_{\mu_1};\,\mathbf{l},\mathbf{l}+\varkappa\Delta_{\mu}}(\Omega_{\mathbf{q}}). \tag{17}$$

Here D is expressed in terms of the Fourier transform of the retarded equal-time Green's function

$$G_{\mathbf{l}_1,\mathbf{l}_1+\varkappa\Delta_{\mu_1};\mathbf{l},\mathbf{l}+\varkappa\Delta_{\mu}}(\Omega_{\mathbf{q}})=-i\int_{-\infty}^{\infty}dt\exp(i\Omega_{\mathbf{q}}t)\theta(t)$$
$$\times\langle[(\tilde{\mathbf{S}}_{\mathbf{l}_1}(t)\tilde{\mathbf{S}}_{\mathbf{l}_1+\varkappa_1\Delta_{\mu_1}}(t)),(\mathbf{S}_{\mathbf{l}}\mathbf{S}_{\mathbf{l}+\varkappa\Delta_{\mu}})]\rangle, \tag{18}$$

where $\langle\cdots\rangle$ denotes averaging over the Gibbs grand ensemble with the Hamiltonian $\mathcal{H}_s$; $\tilde{\mathbf{S}}_{\mathbf{l}}(t)$ is the Heisenberg representation of the spin operator, and $\theta(t)$ is the unit step function:

$$\theta(t)=\begin{cases}1, & t>0\\ 0, & t<0.\end{cases} \tag{19}$$

It is known[7] that the retarded Green's functions are uniquely connected with the Matsubara temperature Green's functions. The latter depend on the discrete frequencies $\omega_n=2\pi nT$, and a well developed diagram-technique formalism has been obtained for them. As applied to spin operators, a diagram technique for Green's functions was developed by Izyumov *et al.*[8] As a result of a calculation of the function (18) in the molecular field approximation, by method developed by them,[8] we get

$$G^{r}_{\mathbf{l}_1,\ \mathbf{l}_1+\nu_1\Delta_{\mu_1};\,\mathbf{l},\,\mathbf{l}+\varkappa\Delta_{\mu}}(\Omega_{\mathbf{q}})=\frac{4b^2(y)n_{\nu}(1+n_{\nu})}{n_{2y}}\frac{\delta_{\mathbf{l}_1,\mathbf{l}+\nu_1\Delta_{\mu_1}}\delta_{\mathbf{l}_1,\mathbf{l}+\nu\Delta_{\mu}}}{2y-\Omega_{\mathbf{q}}-i\delta\,\mathrm{sign}\,\Omega_{\mathbf{q}}} \tag{20}$$

Here $y=g\mu_0H+zJ\langle S^{(3)}\rangle$, z is the number of nearest neighbors ($z=6$ for PCL), $b(y)=SK_s(Sy/T)$, K_s is a Brillouin function, and $n_y=[e^{y/T}-1]^{-1}$.

Substituting (20) and (17), recognizing that the energy absorption D is a real quantity, we obtain after simple transformations

$$D=\frac{8\pi b^2(y)n_y(1+n_y)a^2J^2}{n_{2y}\alpha^2}\frac{1}{N}\sum_{\mathbf{q}\mu}\varepsilon_{\mu\mu}(\mathbf{q})\varepsilon_{\mu\mu}(-\mathbf{q})$$
$$\cos(q_{\mu}a)\,|\Omega_{\mathbf{q}}|\,\delta(2y-\Omega_{\mathbf{q}}). \tag{21}$$

We change over in the usual manner from summation over $\mathbf{q}$ to integration

$$\frac{1}{N}\sum_{\mathbf{q}}\cdots=\frac{a^3}{(2\pi)^3}\int_{-\pi/a}^{\pi/a}dq_1\,dq_2\,dq_3\ldots \tag{22}$$

and write out the integrand in greater detail. Since the dislocation moves along the x_1 axis, we have $\Omega_{\mathbf{q}}=q_1V$. It is easy to show that the products $\varepsilon_{\mu\mu}(\mathbf{q})\varepsilon_{\mu\mu}(-\mathbf{q})$ that enter in (21) take for the linear edge dislocation considered here the form

$$\varepsilon_{11}(\mathbf{q})\varepsilon_{11}(-\mathbf{q})=\frac{8\pi b^2L}{(1-\nu)^2a^6}\frac{q_2^6}{q^8}\delta(q_3), \tag{23}$$

$$\varepsilon_{22}(\mathbf{q})\varepsilon_{22}(-\mathbf{q})=\frac{8\pi b^2L}{(1-\nu)^2a^6}\frac{q_1^4q_2^2}{q^8}\delta(q_3), \tag{24}$$

where ν is the Poisson coefficient; $q=(q_1^2+q_2^2)^{1/2}$; L is the length of the dislocation and coincides with the crystal dimension along the x_3 axis. As a result we obtain from (21) a final expression for the absorption energy per unit dislocation length:

$$\frac{D}{L}=\frac{8b^2(y)n_y(1+n_y)J^2b^2}{n_{2y}\pi(1-\nu)^2\alpha^2a}\theta\left(1-\frac{\zeta}{\pi}\right)\zeta\int_{-\pi}^{\pi}dv\frac{v^6\cos\zeta+\zeta^4v^2\cos v}{(\zeta^2+v^2)^4}. \tag{25}$$

We have introduced here the dimensionless variables $v=q_2a$ and $\zeta=2ya/V$.

3. TEMPERATURE DEPENDENCE OF MAGNETIC DAMPING

Expression (25) contains the function $\theta(1-\zeta/\pi)$, which can be rewritten in the form $\theta(V-2ya/\pi)$. The result should be the inequality

$$V>2ya/\pi, \tag{26}$$

i.e., the velocity has a threshold. If the dislocation moves at less than threshold velocity, the magnetic damping becomes equal to zero under the assumptions made here.

We note that the damping threshold appears always in those cases when the elementary excitations with which the moving dislocation moves has a gap in the spectrum. This pertains, in particular, to spin waves[3] and to electrons in superconductors.[2] In the case considered in the present paper the spins are in an effective molecular field and have equidistant Zeeman energy levels. To excite such a system it is necessary to overcome a threshold corresponding to the first level, i.e., it is necessary that the maximum energy of the Fourier component of the elastic field of the dislocation $\Omega^{\max}=\pi V/a$, transferred to the magnetic system, exceed this threshold value.

We now examine the inequality (26) in greater detail. We introduce for this purpose the value $Y^{\max}$ of y at $T=0$. We consider for simplicity the case $S=1/2$ and neglect the contribution made to y by the external field. Then, within the framework of the molecular-field approximation, we obtain

$$y^{\max}=3J=\Theta_c, \tag{27}$$

$$y=y^{\max}\sigma, \tag{28}$$

where σ is the relative magnetization. Consequently, the inequality (26) can be rewritten in the form

$$V>2\Theta_ca\sigma/\pi. \tag{29}$$

On the other hand, V is limited to the speed of sound $V_s=\Theta_D/k_D$, where Θ_D is the Debye temperature and $k_D\approx\pi/a$ is the Debye momentum. Consequently Θ_c and Θ_D should be connected by the inequality

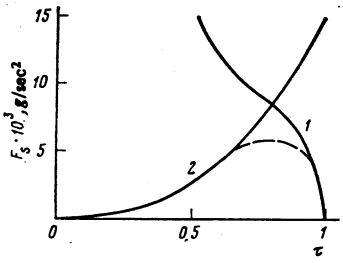

FIG. 1. Temperature dependence of the dislocation damping force. Curve 1 corresponds to magnetic damping in the molecular-field approximation, and curve 2 to damping by spin waves. [3]

$$\Theta_D > 2\Theta_c\sigma. \tag{30}$$

Since σ ranges from unity at $T=0$ to zero at $T=\Theta_c$, it follows that at any ratio of Θ_c and Θ_D the required inequality (30) is satisfied in the region of temperatures close to Θ_c. However, the width of this region depends on the ratio of Θ_c and Θ_D. For classical ferromagnets (Fe, Co, Ni) which have $\Theta_c > \Theta_D$, this temperature region is quite narrow. For ferromagnets[2)] with $\Theta_D > 2\Theta_c$ the width of this region spans the entire temperature interval from zero to Θ_c.

To investigate the temperature dependence of the magnetic damping we assume that the dislocation moves with a velocity V such that the condition (26) is certainly satisfied. We choose for this purpose

$$V = 2ay^{max}. \tag{31}$$

We change over in (25) from energy absorption to the damping force F_s per unit dislocation length and separate the temperature-dependent terms:

$$F_s = \frac{D}{VL} = \frac{Jb^2}{3\pi(1-\nu)^2\alpha^2 a^2} f(T), \tag{32}$$

$$f(T) = \frac{n_\nu(1+n_\nu)}{n_{2\nu}}\sigma^2 \int_{-\pi}^{\pi} dv \frac{v^6\cos\sigma + \sigma^4 v^2\cos v}{(\sigma^2+v^2)^4}. \tag{33}$$

We consider a ferromagnetic material with $\Theta_D > 2\Theta_c$ and choose typical parameter values: $b=a$, $\nu=0.1$, and $\Theta_c = 60$ K. Using the dependence of σ on the temperature for a spin 1/2 in the molecular-field approximation,[9] we obtain from (32) and (33), by computer calculation the values of F as a function of the relative temperature $\tau = T/\Theta_c$ (curve 1 in the figure). From the shape of the curve it is evident that the temperature dependence of F_s in the region of Θ_c is qualitatively correct: the magnetic damping vanishes at $T=\Theta_c$.

For comparison, the figure shows the τ-dependence of the dislocation magnetic damping force[3] calculated by formula (1) with $B_s^{(2)} = 5\times10^{-2}$ g/cm-sec and $V = 3\times10^5$ cm/sec (curve 2). One can thus hope to obtain a qualitatively correct picture of the temperature dependence of the dislocation magnetic damping force in the entire temperature range from zero to Θ_c by joining together curves 1 and 2 in the intermediate temperature region (dashed line). It is seen that at low temperatures the correct dependence is obtained by the spin-wave approximation, where the principal role in the damping is played by magnon scattering by the dislocation phonons,[5] and high energy the correct dependence is given by the molecular-field approximation. In some intermediate temperature region both approaches give results that are close.

4. CONCLUSION

In the preceding sections we investigated the temperature dependence of the magnetic component of the dislocation damping force in a ferromagnet at temperatures of the order of Θ_c. The low-temperature region, which is described in the spin-wave approximation, was investigated by Bar'yakhtar and Druinskiĭ.[3] The intermediate-temperature region can be investigated formally within the approach developed in this paper, but this is fraught with large computational difficulties when it comes to find the Green's function (18) in this temperature region.[8]

A very important question is that of the relative effectiveness of the various dislocation damping mechanisms. For the case $\Theta_D > 2\Theta_c$ considered here, damping of the dislocations by phonons is produced by the flutter-effect mechanism.[1] In this case

$$F_{ph} = B_{ph}(T/\Theta_D)^3 V, \tag{34}$$

where $B_{ph} \approx (10^{-3}-10^{-4})$ g/cm-sec, and the contribution of the phonon damping is negligibly small compared with the magnetic contribution over the entire temperature region from zero to Θ_c, since the phonons are in fact frozen out. Damping by electrons prevails over the magnetic contribution only at very low temperatures, from zero to $T_e \ll \Theta_c$.[3]

As for the classical ferromagnets with $\Theta_c > \Theta_D$, in their case the damping by magnons does not play any significant role: at low temperatures from zero to T_e the main contribution is made by damping by electrons, and at temperatures above T_e, up to Θ_c, it is made by damping by phonons.[3)] Consequently, allowance for the magnetic damping of dislocation is essential only for ferromagnets with $\Theta_c < \Theta_D$. Unfortunately, at present there are no experimental data on dislocation damping in ferromagnets, with which to compare the results obtained on the basis of the theoretical models.

In conclusion, we are sincerely grateful to V. L. Indenbom and A. N. Orlov for a discussion of the work and a number of useful remarks.

1)This is precisely why the calculations are usually confined to the nearest-neighbor approximation.

2)These include, in particular, some rare-earth ferromagnetic elements, such as Er, Ho, and Tu, with low values of Θ_c.

3)It should be noted that Bar'yakhtar and Druinskiĭ[3] obtained an incorrect estimate of the upper limit T_e of the temperature region in which the magnon damping exceeds the phonon damping. A rigorous estimate yields $T_s \sim T_e$, and this excludes the indicated region.

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Multiparticle impurity complexes in uniaxially deformed silicon doped with phosphorus

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The luminescence spectra of multiparticle impurity complexes were investigated at low temperatures in silicon doped with phosphorus and compressed along the [100] direction. It is shown that all the main features of the spectra—the character of line splitting, the line shifts, and the appearance of new lines upon deformation—can be explained on the basis of the shell model of multiparticle complexes.

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Many recent papers are devoted to the physical nature of the series of narrow luminescence lines[1] produced at low temperatures in silicon doped with elements of groups III and V and located in the energy region directly behind the emission line of the bound exciton. According to the interpretation proposed by us earlier[1] these lines are produced when electrons and holes recombine in the multiparticle complexes B_m that are produced as a result of successive binding of m excitons with the impurity atoms. Investigations of the kinetics of the formation and decay of the complexes[2-4] have confirmed this interpretation. However, in subsequent studies[5-7] of the splitting of these lines in magnetic fields and under uniaxial deformation, doubts were cast on the existence of multiparticle complexes. In the present paper, using as an example silicon doped with phosphorus, we show that the complex emission spectrum observed under uniaxial deformation agrees well with the shell model of multiparticle impurity complexes, which was proposed by Kirczenow.[8]

It should be noted that the investigation of the influence of uniaxial compression on the emission spectra of the complexes encounters at least two difficulties. The first is connected with the need for producing a uniform strain that does not cause additional splitting or broadening of the emission lines. In the present study, uniformity of the strain was attained by using long silicon samples measuring ~20×2×2 mm and having hemispherical ends. The pressure P on these ends were applied through lines of a 50% Pb+50% Sn alloy that solidifies at low temperatures.[9] The uniformity of the strain in our was attested by the absence of substantial broadening of the lines up to $P \approx 2000$ kgf/cm^2 when a large area of the crystal with dimension 0.2×7 mm was excited with an argon laser.

The second difficulty arises in the line identification. For reliable identification we used the method of constant additional illumination.[10] To this end we focused on the sample surface, besides the modulated exciting radiation, the beam of a second argon laser. This unmodulated radiation produced a certain stationary concentration of the B_m complexes. As shown in Refs. 1–4, the concentration of complexes with different m depends in nonlinear fashion and in various ways on the excitation levels. Therefore application of constant additional illumination altered the ratio of the amplitudes of the lines of the modulated part of the radiation registered by our apparatus.[10] The fact that the amplitudes of the groups of lines changed in identical fashion was evidence that these lines are connected with one and the same complex B_m.

Figure 1 shows the no-phonon (*NP*) component and the *TA* component of the emission of silicon doped with phosphorus and different pressures in the [100] direction. The emission spectra of the *TA* and *TO* components are similar, and the *TO* spectrum is therefore not shown in Fig. 1. The line designations correspond to the type of radiative transitions assumed in the shell model (Ref. 8). [1)] A number of lines were identified by using constant additional illumination. Thus the onset of the β_2'' line in the decay of the complex B_2 is indicated by the increase of the intensity of this line when the additional illumination is turned on; this increase is identical with that of the α'' line (Fig. 2). The intensity of the line $\alpha_1^{\Gamma_3}$ decreases following the additional illumination in the same way as the intensity of the α_1'' line (Fig. 3), thus indicating that it is produced in the decay of the bound exciton B_1. The spectral positions of the different emission lines of the complexes as functions of the applied pressure P are shown in Fig. 4.

We examine now in greater detail the change of the energy spectrum of multiparticle complexes bound to

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