Dynamic dragging of dislocations

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We analyze, classify, and compare the relative roles of different mechanisms of dislocation dragging, including elastic and inelastic scattering of acoustic and optical phonons by a moving dislocation (phonon wind, Raman scattering, flutter effect) and phonon relaxation (thermoelastic losses, phonon viscosity, relaxation of "slow" phonons). The estimates are carried out with allowance for the finite dimensions of the dislocation core and corresponding deviations of its elastic field from the relations known from the continual dislocation theory. We consider a number of qualitative effects connected with the emission of phonons by a dislocation in the case of stationary motion over the Peierls relief. We discuss the possible mechanisms whereby impurities affect the dynamic dragging of the dislocations. A brief review is presented of the experimental methods of investigating the dynamic mobility of dislocations. The theoretical results are compared with the experimental data on the magnitude and temperature dependence of the dynamic dragging.

CONTENTS

1. INTRODUCTION

A crystal can become plastically deformed because its lattice contains dislocations, which are linear defects capable of being readily displaced under load. When the dislocation is displaced from one surface of the body to another, the crystal experiences a shift along the slip plane by an amount on the order of the lattice parameter a. The motion of the dislocations determines the real atomic structure of a crystal and the kinetics governing the change of the crystal shape under load, and serves as the basis for controlling many important physical properties of solids.

Modern measurement techniques make it possible in principle to trace the displacements of various dislocations, thus providing an experimental base for the investigation of the laws governing dislocation dynamics. The practical importance of the problem in conjunction with the real research possibilities have made dislocation dynamics one of the most rapidly developing branches of solid-state physics.

The study of the mechanisms that limit dislocation mobility under various conditions was initiated approximately a quarter of a century ago in a number of theoretical papers. It had seemed at that time that the main source of dislocation dragging are dissipative processes in the phonon subsystem, which should lead to viscous friction proportional to the dislocation velocity. Subsequent measurements of the mobilities of individual dislocations, however, have revealed that the dragging of the dislocations, generally speaking, does not reduce at all to viscous friction. The first attempts to explain the measured dislocation-mobility curves were not successful, since they started from the incorrect premise that dislocation dragging at all velocities is determined by a single mechanism. It became clear only later that the dislocation mobility is limited by the competition between thermal-fluctuation and dynamic processes, the relative roles of which depend on the dislocation velocity (e.g., [1]). This idea has gained many adherents as it became corroborated by experimental facts.

Figure 1 shows plots of the average slip velocity v of a straight-line dislocation against the applied stress σ, measured in various crystals. An important common property can be traced in the curves plotted in a wide velocity interval. On each of these curves there are two qualitatively distinct stages of the dependence of the velocity on the stress. The first stage is characterized by an abrupt increase of the velocity, by many orders of magnitude, following a relatively small growth of the stress (within one order of magnitude). In the second stage, which begins in the high-velocity region (usually at $v \gtrsim 10^{-2}c$, where c is the speed of sound), the abrupt behavior of the v(σ) curve gives way to a linear dependence of the velocity on the stress, i.e., the dislocation motion

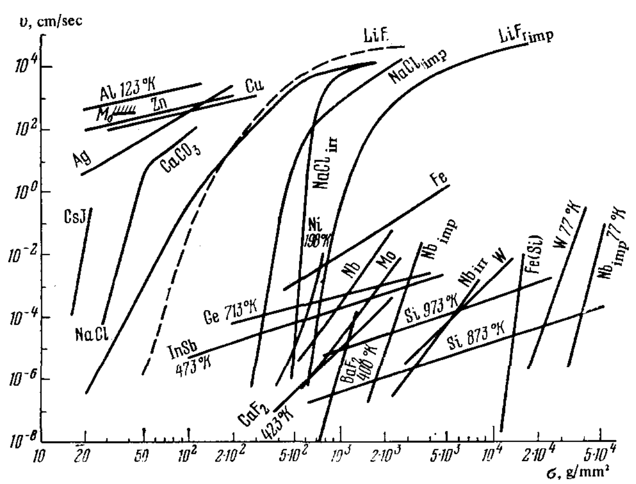


FIG. 1. Mobility curves of individual dislocations for various crystals (see [2]). The subscripts "imp" and "irr" refer to impurity and irradiation, respectively.

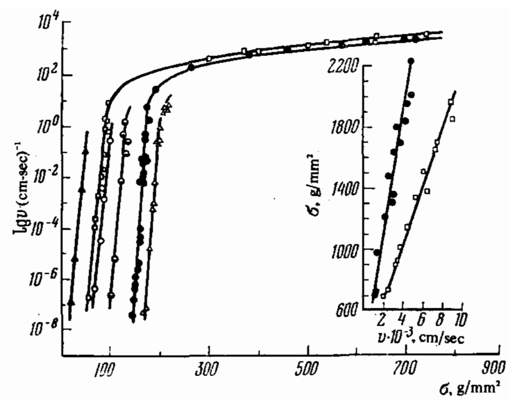
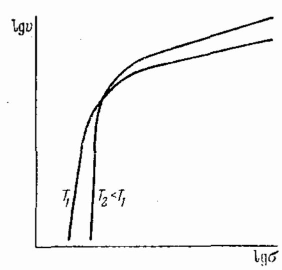


FIG. 2. Dependence of the velocity of individual dislocations on the applied stress in KCl crystals of varying purity (from the data of [3]).

acquires a viscous character. In Fig. 2, this phenomenon is illustrated by plots of the dislocation mobility in KCl crystals with different impurity contents. Figure 3 shows schematically the shift of the v(σ) curve with changing temperature. The experiments indicate without a doubt that there are qualitative differences in the character of the influence of the temperature and of the impurity content on the mobilities of the slow and fast dislocations. Whereas in the first stage small changes of the temperature and impurity concentration lead to an abrupt change of the dislocation mobility, in the second stage the influence of the temperature and of the impurity is much weaker, and the dislocation mobility decreases with increasing temperature at high velocities instead of increasing.

At present one can no longer doubt the validity of the theoretical scheme that attributes the difference in the behavior of the fast and slow dislocations to a radical change in the drag mechanism. For the dislocations to move in the crystals it is necessary that they overcome various types of barriers, which are connected both with the periodic structure of the crystal (Peierls barriers) and with the lattice defects. Slowly moving dislocations stop in front of these barriers and overcome them with thermal fluctuations. The increase of the mobility of the slow dislocations with increasing temperature is due to the higher probability of the thermal fluctuations. With increasing dislocation velocity, when their kinetic energy reaches the height of the energy barriers, conditions are produced for dynamic surmounting the obstacles. The dislocation dragging acquires a dynamic character and is limited by energy transfer from the dislocation to various elementary excitations in the crystal. In contrast to the region of thermal-fluctuation mobility, the dislocation velocity in the dynamic region decreases with temperature because of the increased density of the elementary-excitation gas.

FIG. 3. Effect of temperature on the (σ) plot (schematic).



Dynamic dragging occurs not only for fast dislocations but also for slow ones, and determines the rate of damping of the oscillatory motion of the dislocation segments between the barriers, and even the kinetics of the thermal-fluctuation surmounting of the potential barriers by the dislocations. The first of these effects can be estimated from the amplitude-independent internal friction. The second effect becomes manifest under conditions of abrupt changes in the density of the elementary excitations (for example, if a jumpwise change takes place in the density of the normal electrons as a result of a transition of the metal into the superconducting state).

In most cases, the decisive role in the dynamic dragging of the dislocations is played by dissipative processes in the phonon subsystem of the crystal. The contribution of other types of elementary excitations (electrons, excitons, etc.) becomes manifest only under special conditions. For example, in metals at low temperatures, when the phonon gas is frozen out, an important role is assumed by the interaction of the moving dislocations with the conduction electrons. In this article we confine ourselves to an analysis of the phonon mechanisms of dislocation deceleration.

The first attempt at a comprehensive analysis of the various dissipation channels in the phonon subsystem was made more than 10 years ago in the well known paper of Lothe [4]. At that time, when the experimental investigation of dislocation dynamics had barely started and the role of the dynamic slowing down of the dislocations was not yet fully clear, Lothe's approximate analysis of the contribution of various effects to dislocation drag at room temperatures seemed to be perfectly satisfactory. The later review articles [5,6] and monographs [7,8], containing special chapters on phonon dragging of dislocations, remained in fact within the framework of Lothe's analysis.

The contemporary level of development of the experimental research on dislocation dynamics imposes qualitatively new requirements on the theory. Recently, theoretical papers of more general characters have appeared and have made it possible to establish a definite hierarchy of the phonon mechanisms of dislocation dragging under various conditions, particularly as a function of the temperature. Owing to the latest results of the theory, it becomes possible to explain the bulk of the experimental material on the dynamic mobility of dislocations. There are, however, some questions still to be answered. On the other hand, the theory predicts a number of effects which have not yet been experimentally observed.

The modern ideas concerning the mechanisms of dynamic dragging of dislocations has a history of its own. The verification of a number of physical aspects of dislocation-phonon interaction was accompanied by prolonged discussions, evidence of which can still be seen in the literature. It is necessary to deal with papers in which the formulation of the problem corresponds to a stage of the theory no longer current. We shall attempt below to summarize the present status of the problem, to present a consistent theoretical analysis of the phonon mechanism of dislocation slowing down, to clarify a number of debatable problems, and also to formulate the main problems that still await solution.

2. ANHARMONIC PHONON MECHANISMS OF DRAGGING UNIFORMLY MOVING DISLOCATIONS

The elastic field of a dislocation moving in a crystal disturbs the equilibrium of the phonon gas. As a result, energy flows away from the dislocations to the phonons and the dislocations are effectively slowed down. We shall henceforth consider "nonrelativistic" dislocations, that move with velocities v that are small in comparison with the speed of sound c. In this case the elastic deformation field of the moving dislocation is determined with sufficient accuracy (the correction is of the order of $(v/c)^2$) by the quasistatic transport of the static field at a velocity $\mathbf{v}$: $\epsilon_{ij}(\mathbf{r}, t) \approx \epsilon_{ij}(\mathbf{r} - \mathbf{v}t)$. It is convenient to represent this field in the form of a packet of plane waves, by expanding the function $\epsilon_{ij}(\mathbf{r} - \mathbf{v}t)$ in a Fourier integral:

$$\varepsilon_{ij}(\mathbf{r}, t) = \int \frac{d\mathbf{q}}{(2\pi)^3} \varepsilon_{ij}^{\mathbf{q}} e^{i(\mathbf{qr} - \Omega_{\mathbf{q}} t)}; \qquad (2.1)$$

here $\epsilon_{ij}^{\mathbf{q}}$ is the Fourier transform of the static deformation field of the dislocation, and $\Omega_0 = \mathbf{q}\cdot\mathbf{v}$. Owing to the nonlinear properties (anharmonicity) of the crystal, the phonons should interact with plane elastic waves and this should lead to damping of the packet (2.1). In the linear-response approximation, the energy dissipated per unit time during the dislocation motion is made up of the dampings of the individual waves of the packet (2.1):

$$D = \int \frac{d\mathbf{q}}{(2\pi)^3} \Omega_{\mathbf{q}}^2 \eta_{ijkl}(\mathbf{q}, \Omega_{\mathbf{q}}) \varepsilon_{ij}^{\mathbf{q}} \varepsilon_{kl}^{-\mathbf{q}}, \qquad (2.2)$$

where $\eta_{ijkl}(\mathbf{q}, \Omega_{\mathbf{q}})$ is the effective viscosity for the wave with wave vector $\mathbf{q}$ and frequency $\Omega_{\mathbf{q}}$. The dissipation (2.2) corresponds to the coefficient of dynamic dragging of the dislocations B, which is measured directly in the experiment and is defined as the coefficient of proportionality of the velocity $\mathbf{v}$ and the viscous-dragging force F of the dislocation (per unit length):

$$B = \frac{F}{v} = \frac{D}{v^2}. \qquad (2.3)$$

Typical measured values of B for different crystals at room temperature are 10^{-4}—10^{-3} poise, which corresponds to the viscosity of a rather dense gas.

The coefficients of $\Omega_{\mathbf{q}} \epsilon_{ij}^{\mathbf{q}} \epsilon_{kl}^{-\mathbf{q}}$ in (2.2), which correspond to absorption of individual waves from the packet (2.1), can be interpreted as the imaginary parts of the dynamic elastic moduli $c_{ijkl}(\mathbf{q}, \Omega_{\mathbf{q}})$. In terms of these concepts, Kosevich and Natsik [9] have developed a phenomenological theory of dislocation dragging, the results of which have a particularly elegant form in the case when spatial dispersion c_{ijkl} is neglected. This theory turned out, in particular, to be fruitful in the analysis of the role of quasilocal oscillations of impurity centers excited when the dislocation moves (see Sec. d of Chap. 3). Unfortunately, the problem of phonon dragging of dislocations is made complicated by the fact that the predominant contribution to the dissipation is made by processes having a sharply pronounced spatial dispersion.

When estimating the dislocation dragging it must be taken into account that, depending on the scale of the temporal and spatial inhomogeneities of the external field, the dissipative processes in the phonon subsystem are qualitatively different. As applied to the analysis of the damping of a plane wave, the natural quantities with which to compare its frequency ω and its wavelength $\lambda = 2\pi/q$ are the reciprocal phonon relaxation time τ^{-1} and the mean free path $l = c\tau$. In Fig. 4a, the phase plane $\{\omega, q\}$ is subdivided into four regions characterized by the parameters $\omega\tau$ and ql. We are interested in only three regions in which the frequencies $\Omega_{\mathbf{q}}$ and wave vectors $\mathbf{q}$ of the plane waves from the packet (2.1) can fall. It is clear from general considerations that in region I ($\omega\tau < 1$, $ql < 1$) the dissipative processes have a macroscopic and relaxation character, in the region II ($\omega\tau > 1$, $ql > 1$) quantum processes of phonon scattering predominate, and finally in the region III ($\omega\tau < 1$, $ql > 1$) the dissipative processes should include both scattering and relaxation. We shall analyze below the corresponding estimates of the phonon dragging in different macroscopic and microscopic approximations, and then formulate a general approach that admits of a unified description of the dissipative processes in all regions I—III.

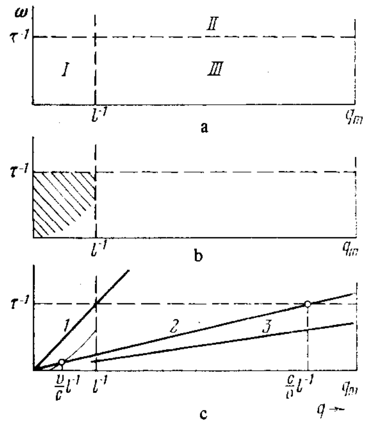


FIG. 4. Phase diagram illustrating the different behavior of the damping of a plane elastic wave interacting with the phonon sub-system of a crystal, depending on the region of the $\{\omega, q\}$ plane in which its frequency and wave vector fall. I—Region where macroscopic relaxation processes such as phonon viscosity and thermoelastic losses predominate (there is no spatial dispersion of the damping in the shaded parts of region I ($\omega > \chi q^2$)); II—quantum region, in which the principal role is played by phonon scattering by the elastic field of the wave; III—kinetic region, where the scattering and relaxation of the phonons compete. Line 1—dispersion law $\omega = cq$ for sound waves; lines 2 and 3—dispersion law $\omega = vq$ for the "fast" ($v > cq_m l$) and "slow" ($v < cq_m l$) dislocation packets (2.1), respectively.

a) Phonon viscosity. The damping of plane elastic waves from the packet (2.1) is obviously limited by dissipative processes of the same type as absorption of ultrasound. Akhiezer has shown at one time [10], in an analysis of the absorption of long-wave ultrasound, that in the course of restoration of the equilibrium of a phonon gas disturbed by a sound wave, the phonons behave like a gas with effective viscosity $\eta\gamma^2 E\tau$ (E is the thermal-energy density and γ is the Gruneisen constant). It is therefore natural to attempt to estimate the dislocation dragging by substituting $\eta(\mathbf{q}, \Omega_{\mathbf{q}}) = \eta$ in (2.2) (we omit the tensor indices for now on for the sake of simplicity):

$$D = \eta \int \frac{d\mathbf{q}}{(2\pi)^3} \Omega_{\mathbf{q}}^2 |\varepsilon^{\mathbf{q}}|^2. \qquad (2.4)$$

An estimate of just this type, obtained by Mason [11], is still being used by some workers in the analysis of experimental material. Since the deformations near the dislocation increase in inverse proportion to the distance, the integral (2.4) diverges quadratically at the upper limit and should be cut off at a certain value q_m

$\sim r_0^{-1}$. In his calculations, Mason started from the fact that r_0 has the meaning of the radius of the dislocation core, and therefore should be of the order of the lattice parameter a.

It is easy to verify, however, that the estimate (2.4) is in principle incorrect and is even of the wrong order of magnitude. As already noted, the main contribution to the integral of (2.4) is made at large q, but it is precisely at large q that the neglect of the spatial dispersion of the phonon viscosity in (2.4) is inadmissible and leads to incorrect results. The concept of phonon viscosity without dispersion can be introduced only in the "adiabatic" part of region I, which is shown shaded in Fig. 4b, and for which

$$\omega \gg \chi q^2, \tag{2.5}$$

where $\chi \approx cl/3$ is the thermal diffusivity (for details see Sec. e of this chapter, the second term of formula (2.26)). It is seen from Fig. 4c that in contrast to the acoustic wave, which is adiabatic in the entire region I, the partial waves from the packet (2.1) are adiabatic only in a small region $q < (v/c)l^{-1}$. Thus, from formula (2.4) we can estimate the damping of only the long-wave part of the packet (2.1), which makes a negligibly small contribution to the dislocation dragging: the use of the quantity $(c/v)l$ as the cutoff radius in place of a decreases Mason's estimates by $(cl/va)^2$ times!

In spite of all the foregoing, Mason's paper [11] played a positive role in the development of the modern ideas concerning dissipative processes that determine the dynamic dragging of the dislocations. First, the idea of the fundamental difference between the mechanisms that limit the mobilities of the fast and slow dislocations seems to have been advanced for the first time in his paper. In addition, the subsequently observed fair agreement between Mason's formula and certain experimental data on the temperature dependence of the dynamic slowing down of dislocations stimulated further searches for a relaxation mechanism that leads to a dragging proportional to τ. Indeed, as we shall see, the concept of a constant phonon viscosity, which is not applicable to ordinary "acoustic" phonons, turns out to be applicable to "slow" phonons with low group velocities. The relaxation of these "slow" phonons will be considered in Secs. e and f of the present chapter.

b) Thermoelastic dissipation. Dislocation motion leads to heating and cooling of the crystal sections subjected to rapid compression and tension. This causes additional fluxes of phonons—heat flow from a hot section to a cold one, accompanied by thermoelastic energy dissipation [12]:

$$D = \frac{\varkappa}{T}\int (\nabla T)^2\, d\mathbf{r} = \frac{\varkappa}{T}\int \frac{d\mathbf{q}}{(2\pi)^3}\,|(\nabla T)_{\mathbf{q}}|^2; \tag{2.6}$$

here $\kappa = C\chi$ is the thermal-conductivity coefficient, C is the specific heat of the crystal, T is the average temperature of the crystal, and ∇T and $(\nabla T)_q$ are the temperature gradient and its Fourier transform. The thermoelastic dissipation can be treated as a particular case of phonon viscosity, wherein the deformation reduces to dilatation, the change in the phonon spectrum reduces to local heating (or cooling of the crystal), and restoration of the equilibrium in the phonon subsystem reduces to heat flows. The heat-conduction equation in a crystal with moving dislocation is of the form

$$\frac{\partial T}{\partial t} - \chi\Delta T = \gamma T \dot{\varepsilon}_{ll}(\mathbf{r}, t); \tag{2.7}$$

the right-hand side corresponds here to the volume heat release under adiabatic change of the crystal volume. Taking the Fourier transform of (2.7), we easily obtain

$$(\nabla T)_{\mathbf{q}} = \frac{\gamma T \varepsilon^{\mathbf{q}}_{ll}\Omega_{\mathbf{q}} q}{-i\Omega_{\mathbf{q}} + \chi q^2}. \tag{2.8}$$

Let us estimate with the aid of formulas (2.6) and (2.8) the deceleration of a straight-line edge dislocation with a Burgers vector b, moving in a glide plane with a unit normal vector **n**. The continual theory of dislocations yields in the isotropic approximation, as is well known [12],

$$\varepsilon_{ll}(\mathbf{r}) \approx -\frac{b}{4\pi}\frac{(\rho\mathbf{n})}{\rho^2}, \quad |\varepsilon^{\mathbf{q}}_{ll}|^2 \approx \frac{\pi}{2} b^2 \frac{\delta(\mathbf{qm})}{q^2}\left(\frac{\mathbf{qn}}{q}\right)^2, \tag{2.9}$$

where $\boldsymbol{\rho}$ is that component of the radius vector **r** which is perpendicular to the dislocation, **m** is a unit vector directed along this location, $\delta(x)$ is the Dirac delta function, and the Poisson coefficient has been set equal to 1/3 by way of estimate. Taking (2.8) and (2.9) into account, the integration in (2.6) can be carried out directly, and because of the logarithmic divergence with respect to q the integration at the upper limit should be cut off at $q = q_m$

$$D \approx \frac{\gamma^2 b^2}{64\pi} v^2 \frac{CT}{\chi} \ln \frac{\chi q_m}{v}. \tag{2.10}$$

The quantity q_m should be chosen from physical considerations. Estimating the thermoelastic damping, Eshelby [13] (for vibrating dislocation segments) and Weiner [14] (for a dislocation moving in translation) extended the integration in (2.6) all the way to the core of the dislocation, corresponding to the choice $q_m \sim 1/a$. It must be borne in mind, however, that the macroscopic approach developed here is valid only for volume elements with linear dimensions that exceed the phonon mean free path l. In this sense, the value $q_m = \pi/l$ proposed by Lothe [4] seems more natural. The corresponding value of the parameter under the logarithm sign is still large in this case, $\chi q_m/v \approx c/v \gg 1$. Changing over from the dissipation D to the dragging coefficient $B = D/v^2$, we have

$$B = \frac{\gamma^2 b^2}{64\pi}\frac{CT}{\chi}\ln\frac{c}{v}. \tag{2.11}$$

A numerical estimate of the dragging coefficient in accordance with formula (2.11) leads to values that are smaller by 1—2 orders of magnitude than the experimentally observed values. This is connected in part with the fact that the Gruneisen constant is too rough a measure of the anharmonicity of the crystal in similar problems. We shall return later on (Sec. e of this chapter) to a discussion of the relative role of the temperature losses, on the basis of a more general approach to the problem (in particular, with a more complete allowance for the anharmonicity). We note only that the considered mechanism leads to a deceleration of only edge or mixed dislocations, the deformation field of which contains a dilatation component.

Of course, phonon relaxation does not reduce merely to thermoelastic processes, and should take place also in a shear field. We have seen in Sec. a that at distances $R > (c/v)l$ from the dislocation the relaxation can be described in the language of ordinary phonon viscosity. At shorter but still macroscopic distances, $l < R < (c/v)l$, the estimates of the phonon relaxation calls for taking into account the spatial dispersion of the phonon viscosity, and will be considered in Sec. e of this chapter. An analysis of the dissipative processes near the dislocation

(R < l) cannot be carried out in macroscopic terms, and calls for a microscopic quantum approach.

c) Phonon scattering. A description of the phonon processes that occur arbitrarily close to the moving dislocations was proposed by Leibfried[15], who has considered the scattering of phonons by dislocations. In a coordinate system connected with the moving dislocation, the flux of phonons incident on a dislocation is asymmetrical. Owing to this purely aberrational effect, the phonons scattered by the dislocation impart to the latter a momentum proportional to the dislocation velocity and directly opposite to the dislocation motion. The dislocation experiences, as it were, a phonon wind. It is easy to visualize that the momentum F transferred in this case from the phonons to each unit of dislocation length per unit time should be of the order $fd_{tr}E(v/c)$, where d_{tr} is the transport diameter for the scattering of the phonons by the dislocation, and f is the numerical factor ($f \lesssim 1$). The corresponding estimates for the dragging coefficient $B = F/v$ can be usually written down, following Leibfried[15], in the form

$$B = \frac{bE}{10c}. \tag{2.12}$$

Leibfried did not calculate the cross section for phonon scattering by the dislocation, assuming it to be of the order of the Burgers vector b (usually $b \sim a$). In this sense, formula (2.12) applies equally well to any scattering mechanism with a diameter on the order of a. As shown by Nabarro[16], a distinction should be made between two mechanisms of phonon scattering by a dislocation, namely the phonon wind (nonlinear mechanism) due to nonlinear elastic properties (anharmonicity) of the crystal, and the flutter effect, which is connected with reradiation of phonons by a dislocation that oscillates in the thermal field of the lattice. Although Leibfried had in mind scattering due to nonlinearity of the properties of the medium near the dislocation, Lothe has shown subsequently[17] that the estimate (2.12) is valid also for the flutter effect at temperatures on the order of the Debye temperature. Indeed, according to[18] the phonon-scattering diameter due to the nonlinear mechanism is of the order of $\gamma^2 b^2/\lambda$ (λ is the phonon wavelength), while the flutter diameter is of the order of λ. It is therefore understandable that formula (2.12) describes both effects at sufficiently high temperatures, when phonons with wavelength on the order of the lattice parameter predominate. But it is also obvious from this that these mechanisms have different temperature dependences, and the occasionally employed extrapolation of formula (2.12) to the region of low temperatures[11,19,20] is meaningless.

Numerical estimates based on formula (2.12) leads to a dragging coefficient approximately one order of magnitude lower than that observed in experiment. This is due, in particular, to the fact that the Gruneisen constant, as already noted above, is a poor measure of the anharmonicity when it comes to phonon scattering (see Sec. d of this chapter). To obtain reliable information on the magnitude, temperature dependence, and relative role of the two aforementioned scattering mechanisms it is necessary to analyze the problem theoretically in greater detail. We shall discuss the quantum-mechanical theory of phonon wind and of the flutter effect in Sec. d of this chapter and in Sec. a of Chap. 3, respectively.

d) Phonon wind (nonlinear mechanism). We can attempt to estimate the interaction of the phonons with the internal-stress field in the crystal, an interaction due to the anharmonicity of the lattice by means of the dilatation effect that leads to a local change in the phonon velocity $\delta c \approx \gamma c \epsilon_{ll}$ and to a corresponding frequency shift at[18]

$$\omega' = \omega\left(1 + \gamma\varepsilon_{ll}(\mathbf{r}, t)\right). \tag{2.13}$$

The energy of each phonon in the dislocation field is perturbed in this case by an amount $\hbar\omega\gamma\epsilon_{ll}$ ($\hbar$ is Planck's constant). This is precisely the idea on the basis of which the first calculations[21,22] were made of the cross section for phonon scattering by dislocations. However, the estimate of the phonon dragging of the dislocations based on this approach[4] leads to a result of the type (2.12), which is smaller by one order of magnitude than the experimentally observed values. In addition, in the approximation (2.13) there is no dragging (in contradiction to the experiment) of the screw dislocations, for which $\epsilon_{ll} = 0$. It appears that to estimate this effect with any degree of reliability it is necessary to take into account the anharmonicity of the lattice in terms of elastic constants of higher order. This approach to the problem was first used in the papers of Al'shitz[23] and Gruner[24],[1)].

The first anharmonic term in the expansion of the elastic energy of a crystal in terms of the strains can be represented in the continual approximation in the form[18]

$$\Phi_{\text{anh}} = \frac{1}{3!}\int d\mathbf{r} A^{lnj}_{kmi} u_{kl} u_{mn} u_{ij}; \tag{2.14}$$

here A^{lnj}_{kmi} is the third-order elastic-constant tensor. In a crystal with a dislocation it is necessary to add to the thermal strains of the crystals u_{ij} the dislocation strains ϵ_{ij}. As a result, in the approximation linear in the dislocation field, the energy of interaction of the thermal oscillations with the dislocation takes the form

$$\Phi_{\text{int}} = \frac{1}{2}\int d\mathbf{r} A^{lnj}_{kmi} u_{kl} u_{mn} \varepsilon_{ij}, \tag{2.15}$$

which leads, after the usual transition to Fourier space (with allowance for (2.1)) and second quantization in terms of the phonon variables, to the following Hamiltonian of the interaction between the phonons and the moving dislocation[26]:

$$\mathscr{H}_{\text{int}}(t) = \sum_{\alpha,\beta} \Gamma_{\alpha\beta}\xi^{+}_{\alpha}\xi_{\beta} e^{-i\Omega_{\mathbf{q}}t}; \tag{2.16}$$

the subscripts α and β denote here different states of the phonons, specified by the aggregate of the wave vector $\mathbf{k}$ and the polarization λ: $\alpha = (\mathbf{k}, \lambda)$, $\beta = (\mathbf{k}', \lambda')$, with $\mathbf{k}' = \mathbf{k} + \mathbf{q}$; $\xi_\alpha = a_{\mathbf{k}\lambda} + a^{+}_{-\mathbf{k}\lambda}$, $a^{+}_{\mathbf{k}\lambda}$ and $a_{\mathbf{k}\lambda}$ are the phonon creation and annihilation operators:

$$\Gamma_{\alpha\beta} = A^{ij}_{\alpha\beta}\varepsilon^{\mathbf{q}}_{ij}, \quad A_{\alpha\beta} \sim \frac{M}{4G}\frac{\hbar c^2 kk'}{\sqrt{\omega_\alpha\omega_\beta}}, \tag{2.17}$$

M is the characteristic value of the third-order modulus and G is the shear modulus. Since an energy $\hbar\Omega_{\mathbf{q}}$ is transferred in each act of phonon scattering from the state α to the state β and back, it follows that the energy dissipation per unit time is

$$D = \frac{\pi}{\hbar^2}\sum_{\alpha,\beta}\hbar\Omega_{\mathbf{q}}|\Gamma_{\alpha\beta}|^2(n_\beta - n_\alpha)\,\delta(\omega_\alpha - \omega_\beta - \Omega_{\mathbf{q}}) \approx \approx -\frac{\pi}{\hbar}\sum_{\alpha,\beta}\Omega^2_{\mathbf{q}}|\Gamma_{\alpha\beta}|^2\frac{\partial n_\alpha}{\partial\omega_\alpha}\delta(\omega_\alpha - \omega_\beta); \tag{2.18}$$

Here $n_\alpha = [\exp(\hbar\omega_\alpha/k_B T) - 1]^{-1}$ is the equilibrium phonon distribution function and k_B is the Boltzmann constant.

Expression (2.18) describes the dragging of not only a dislocation but also of any elastic-field source (for example a kink on a dislocation, a crowdion, etc.), provided that the tensor ϵ_{ij}^{q} is suitably chosen, and naturally coincides with the analogous formula [27] which deals with the analysis of the nonlinear mechanism of phonon dragging of a kink.

Substituting expression (2.17) in (2.18) and recognizing at the same time that $|\epsilon_{ij}^{q}|^2$ is determined in order of magnitude by formula (2.9), we obtain, after replacing summation by integration and after performing a number of simple calculations, the required estimate of the contribution of the phonon wind to the dislocation dragging:

$$B=\frac{D}{v^2}=\left|\frac{M}{G}\right|^2\frac{\hbar}{b^3}\left(\frac{k_D b}{2\pi}\right)^5 f\left(\frac{T}{\Theta}\right); \qquad (2.19)$$

here

$$f(x)=x^5\int_0^{1/x}\frac{e^t t^5\,dt}{(e^t-1)^2}, \qquad (2.20)$$

Θ is the Debye temperature, and k_D is the Debye end point in the phonon spectrum.

It follows in particular from (2.19) and (2.20) that at low temperatures we have $B \propto T^5$, whereas extrapolation of (2.12) to this region yields $B \propto T^4$.

The coefficient $|M/G|^2$ for a screw dislocation was determined in [23] in the isotropic approximation:

$$\left|\frac{M}{G}\right|^2\approx 4+\left(\frac{|n|}{G}-6\right)^2, \qquad (2.21)$$

where n is the Murnaghan modulus. Usually $|n|/G \sim 15-30$, and therefore $|M/G|^2 \sim 10^2-10^3$. In [24] there was obtained a low-temperature asymptotic expression for the dragging coefficient B for screw and edge dislocations in copper. The coefficient of T^5, as obtained by Gruner, is approximately seven times larger than the coefficient corresponding to the estimate of (2.19) and (2.21). It is impossible to explain the cause of the discrepancy, since only the final result is presented in [24], without proof, and furthermore in a rather strange form. The values expressed in terms of the energy density of the longitudinal phonons and the Debye temperature, whereas the main contribution to the slowing down is made by the transverse phonons[2)], and the low-temperature asymptotic form of (2.19) does not depend on Θ.

At high temperatures, formula (2.19) is much less reliable, for when the average phonon wavelength is decreased the problem becomes more and more sensitive to the deviation of the phonon spectrum from the Debye spectrum, and to the structure of the deformation field in the core of the dislocation. Therefore, to obtain the temperature dependence of B in the entire region it is necessary, strictly speaking, to know the true phonon spectrum of the crystal, and also the dislocation field in the region where the continual theory does not hold. In principle, the experimentally obtained temperature dependence of B(T) can be used to obtain information on the phonon spectrum and on the structure of the dislocation core. We confine ourselves here to a simple illustration, which shows how allowance for the existence of a dislocation core influences the temperature dependence of B(T). We introduce as the simplest model of the core a smooth cutoff of the dislocation field at short distances:

$$\varepsilon_{ij}(\mathbf{r})=\varepsilon_{ij}^0(\mathbf{r})\,(1-e^{-r/r_0}); \qquad (2.22)$$

Here $\epsilon_{ij}^0(\mathbf{r})$ is the strain tensor in the continual-theory approximation and r_0 is the effective radius of the dislocation in the nucleus (according to [4], $r_0 \approx 3b$). The Fourier transform of the field $\epsilon_{ij}(\mathbf{r})$ is also altered, and to estimate $|\epsilon_{ij}^{q}|^2$ it is necessary to use in place of (2.9) the expression

$$|\varepsilon_{ij}^q|^2=\frac{\pi}{2}\frac{\delta(\mathbf{qm})}{q^2}\frac{b^2F_{ij}(\mathbf{q}/q)}{1+(r_0q)^2}, \qquad (2.23)$$

where $F_{ij}(\mathbf{q}/2)$ is a function of the directions and is of the order of unity. Formula (2.20), which describes the temperature dependence of the effects, should accordingly be replaced by

$$f_1(x)=x^5\int_0^{1/x}\frac{dt\,e^t t^5}{(e^t-1)^2}\frac{\operatorname{arctg}\beta xt}{\beta xt}, \qquad (2.24)$$

where $\beta = 2k_D r_0$. At $\beta \ll 1$, the function $f_1(x)$ goes over naturally into f(x). It is easy to verify, however, that usually we have, to the contrary, $\beta \gg 1$ ($\beta \sim 30$ for typical values of k_D and $r_0 \approx 3b$). It is seen from (2.24) that at low temperatures ($\beta x \ll 1$) the function $f_1(x)$ is practically independent of β and coincides with f(x). With increasing temperature, the function $f_1(x)$ assumes quite rapidly the linear form

$$f_1(x)\approx\frac{x}{3\beta}\left[\operatorname{arctg}\beta-\frac{\beta^2-\ln(\beta^2+1)}{2\beta^3}\right]\approx\begin{cases} x/4, & \beta\ll 1,\\ x/2\beta, & \beta\gg 1.\end{cases} \qquad (2.25)$$

Since usually $\beta \gg 1$, as we have seen, allowance for the finite dimensions of the dislocation nucleus decreases the estimate (2.19) and (2.20) at high temperatures by a factor $\beta/2 = k_D r_0$, i.e. by approximately one order of magnitude. This eliminates the contradiction between theory and experiment, which was noted by Brailsford [25], who estimated the dragging coefficient for copper at room temperature by means of formulas similar to (2.19) and (2.20) and obtained a value of B exceeding the measured values by more than one order. We shall show below that at room temperature formulas (2.19) and (2.24) lead to values of B close to those observed in experiment, unlike the overestimates that do not take into account the relaxation of the elastic dislocation field near the core, and unlike the underestimates in which the Gruneisen constant γ is used as a measure of the anharmonicity.

The reason for the discrepancy between the absolute values of the estimates based on different methods of anharmonicity was explained in a paper by Al'shitz [23]. Expressing the Gruneisen constant in terms of the Murnaghan moduli, the author has shown that the modulus n, which determines, as we have seen, the amplitude of the effect makes a relatively small contribution to the constant γ, and furthermore causes it to decrease. For this reason, a direct comparison of the estimates (2.11) and (2.19), which were obtained in different approximations, is incorrect.

The foregoing calculation of the phonon wind does not take into account the relaxation of the phonon gas in the process of scattering by a dislocation; this neglect is permissible only in the case $\Omega_q \tau \gg 1$, i.e., for the short-wave part of the packet (2.1) ($q \gg (c/v)l^{-1}$), which belongs to region II on the $\{\omega, q\}$ phase diagram (see Fig. 4). It is therefore necessary to introduce a cutoff from below in (2.18). This changes the estimate (2.19) little at sufficiently rapid dislocations and low temperatures, but shows by the same token that the calculation cannot be applied to slow dislocations and high temperatures, when $(c/v)l^{-1} > q_m$, and the phase line $\omega = qv$ does not fall in the region II at all (line 3 in Fig. 4c).

Thus, the methods considered above make it possible to investigate only processes very close to the dislocation or very far from it. An estimate of the contribution made to the dissipation of the intermediate region $(v/c)l < R < (c/v)l$, where the calculation should take into account both the scattering and the relaxation of the phonons, calls for a solution of the kinetic problem for the phonons in the field of the moving dislocation. The kinetic approach, as will be shown below, has the advantage that it makes it possible to consider from a unified point of view different dissipation channels and to determine their relative contribution to the dragging of the dislocation.

e) Unified approach to the scattering and relaxation processes. To describe different phonon mechanisms of dislocations dragging within the framework of a single formalism one can use, in principle, the classic kinetic equation, as proposed by Brailsford[25]. We regard as more consistent, however, the quantum approach developed by Al'shitz and Mal'shukov[26], in which it is possible to monitor the errors in the derivation of the kinetic equation, and if necessary to go outside the framework of the usual Boltzmann equation. Incidentally, the analysis in[25,26] contains, as we shall show, one common significant gap, namely, the use of the Debye model of the phonon spectrum has made it impossible to note and to analyze one more important dissipation channel, which no one has noted before, namely the relaxation of "slow" phonons. We shall attempt below to fill this gap.

The perturbation of the phonon gas in the field of a moving dislocation should lead to a deviation $\Delta\rho$ of the phonon density matrix from the equilibrium value ρ_0. In the first-order approximation in the perturbation (2.16), the quantity $\Delta\rho$ is linear in $\mathcal{H}_{\text{int}}$, and the dissipation of the energy per unit time $D = -\text{Sp}\,(\Delta\rho\,\partial\mathcal{H}_{\text{int}}/\partial t)$ is correspondingly quadratic in $\mathcal{H}_{\text{int}}$. It can be shown[26] that the dissipation D is determined by formula (2.2), in which the effective viscosity $\eta_{ijkl}(\mathbf{q}, \Omega_{\mathbf{q}})$ is the Fourier transform of the two-particle retarded Green's function for the phonons. Calculation of this function in the harmonic approximation leads, naturally, to an expression for D identical with (2.18). The relaxation processes can be taken into account by expanding the Green's function in a perturbation-theory series in the anharmonicity. It is convenient to use the well-developed diagram technique.

In the long-wave region ($ql \ll 1$), the perturbation-theory series contains singular ladder diagrams, the summation of which is equivalent to the solution of the kinetic equation. Leaving out the cumbersome intermediate steps, we present only the final results of the realization of the procedure described above. The effective viscosity $\eta_{ijkl}(\mathbf{q}, \Omega_{\mathbf{q}})$ in the region $ql \ll 1$ is determined by the expression

$$\eta_{ijkl}(\mathbf{q}, \Omega_{\mathbf{q}}) = \frac{\alpha^2 CT\tau}{3(1+p^2)}\left[\left(\frac{\omega_q}{\Omega_q}\right)^2 \delta_{ij}\delta_{kl} + \varphi_{ijkl}\left(\frac{\mathbf{q}}{q}\right)\right]; \qquad (2.26)$$

here $p = \chi q^2/\Omega_{\mathbf{q}}$ is the adiabaticity parameter, $\omega_{\mathbf{q}} = qc$, δ_{ij} is the Kronecker symbol, $\varphi_{ijkl}(\mathbf{q}/q)$ is an angle function of the order of unity, and α is the numerical coefficient proportional to the anharmonicity level of the crystal. The values of α used in[25,26] were respectively $\alpha_B = \gamma$ and $\alpha_A = M/4G$. The estimates expressed in[25] in terms of the Gruneisen constant γ do not claim to yield the correct absolute values, but nevertheless make it possible to assess the relative roles of processes of different types. To estimate the absolute values of the investigated effects, Brailsford proposed a phenomenological approach based on the normalization of the dislocation interaction against the experimental data on the influence of the dislocations on the thermal conductivity. It appears that this method is not too reliable, so that at present there is no unambiguous theoretical interpretation of the existing data on the dislocation component of the thermal conductivity. In particular, it is thought (see the discussion at the conference[28]) that the effect is determined by the flutter mechanism, which is not connected at all with the anharmonicity of the crystal.

The first term of (2.26) corresponds to thermoelastic damping, and the second to phonon viscosity. In ultrasound-absorption theory one uses the fact that at $ql \ll 1$ the sound wave is adiabatic: $p \ll 1$ (see formula (2.5)), so that the dispersion of the phonon viscosity and of the thermoelastic losses can be neglected. It is easily seen, however, that when applied to the analysis of the damping of the dislocation packet (2.1), this neglect is permissible only for the longest wavelength: $q \ll (v/c)l^{-1}$, and it is on this fact that our criticism of Mason's theory, given in Sec. (a), is based. It is seen from (2.26) that independently of the spatial dispersion the contribution made to the dragging of the dislocation by the phonon viscosity is smaller by a factor $(c/v)^2$ than the thermoelastic losses, in contrast to the case of ultrasound, when $\Omega_{\mathbf{q}} \approx \omega_{\mathbf{q}}$ and both terms in (2.26) are of the same order. We shall neglect henceforth the second term in (2.26) in comparison with the first.

At $ql > 1$, the diagrams do not contain any singularities. This enables us to use the relaxation-time approximation in the calculation. As a result, the damping of the short-wave part ($ql \gg 1$) of the packet (2.1) is given, accurate to terms proportional to l^{-1}, by the formula[26]

$$D = -\frac{\pi}{\hbar}\sum_{\alpha,\beta}\Omega_{\mathbf{q}}^2\,|\Gamma_{\alpha\beta}|^2\,\frac{\partial n_\alpha}{\partial\omega_\alpha}\left[\frac{(\pi\tau)^{-1}}{(\omega_\alpha-\omega_\beta)^2+\tau^{-2}}\right]. \qquad (2.27)$$

The expression in the square brackets, which we designate for convenience by $Q_{\alpha\beta}(\tau)$, goes over formally as $\tau \to \infty$ into $\delta(\omega_\alpha - \omega_\beta)$, and in this case formula (2.27) coincides with (2.18). The degree to which $Q_{\alpha\beta}(\tau)$ approaches a delta function depends not only on the value of the parameter $\omega_\alpha\tau$, but also on the character of the function $\omega_\alpha = \omega(\mathbf{k})$, since the integration in (2.17) is carried with respect to $\mathbf{k}$ and not with respect to ω. The estimates in[25,26] were obtained using a formula similar to (2.27) in the Debye approximation, when $\omega_\alpha(\mathbf{k}) = ck$ and we have with good accuracy $Q_{\alpha\beta}(\tau) \approx \delta(\omega_\alpha - \omega_\beta)$. It is therefore natural that the correspondingly expression for η_{ijkl} is independent of the relaxation time τ:

$$\eta_{ijkl}(\mathbf{q}, \Omega_{\mathbf{q}}) \sim -\left(\frac{\alpha}{2}\right)^2\frac{\hbar c}{q}\int_{q/2}^{k_D} dk\,k^4\,\frac{\partial n}{\partial\omega}. \qquad (2.28)$$

Expression (2.28) is reliable enough only at low temperatures, when long-wave phonons predominate. As already noted above, at high temperatures it becomes necessary to take into account the deviation of the real phonon spectra from the Debye model. Since the dispersion of the short-wave phonons deviates from linearity (see the plot in Fig. 5), the fraction of the "slow" phonons with small group velocities $v_{gr} = \partial\omega/\partial\mathbf{k}$ increases with increasing temperature. In particular, near the boundary of the Brillouin zone there is a region in which the phonon dispersion is smaller than the natural line width, $|\omega_\alpha - \omega_\beta| \ll \tau^{-1}$. The contribution made to the phonon dissipation

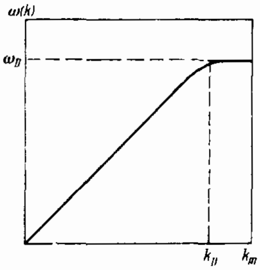


FIG. 5. Typical form of dispersion curve ω(k) for acoustic phonons.

by this region differs qualitatively from the contribution of the "acoustic" phonons, inasmuch as for the "slow" phonons we have $Q_{\alpha\beta}(\tau) \approx \tau/\pi$. The requirement $|\omega_\alpha - \omega_\beta| \ll \tau^{-1}$, of course, does not necessarily mean that the phonon group velocity v_{gr} is literally small. It is only necessary that the energy of the phonon in the final state with wave vector **k** + **q** differ little from its energy in the initial states **k**. Taking into account the orthogonality of the vectors **q** to the dislocation line (see formulas (2.9) and (2.23)), we can easily verify that this property is possessed, for example, by phonons belonging to the flat sections of the equal-frequency surfaces, when these flat sections are perpendicular to the dislocation line. An analysis of real phonon spectra shows that such flat sections, oriented parallel to the boundary of the Brillouin zone, usually occupy an appreciable volume of phase space.

We confine ourselves below to a simplified estimate of the role of the "slow" phonons, assuming that the region $|\omega_\alpha - \omega_\beta| \ll \tau^{-1}$ occupies a finite phase volume amounting to a fraction g of the volume of the Brillouin zone, and is strongly enough pronounced to be able to neglect the contribution of the transition region. Substituting $Q_{\alpha\beta}(\tau) = \tau/\pi$ in (2.27), we can easily estimate the contribution made to (2.28) by the relaxation of the "slow" phonons:

$$\eta_{ijkl}(\mathbf{q}, \Omega_q) \sim \alpha^2 C_S T \tau \Psi_{ijkl}(\mathbf{q}); \qquad (2.29)$$

here C_S is the contribution of the "slow" phonons to the specific heat of the crystal, $\Psi_{ijkl}(\mathbf{q})$ is a function of the order of unity and decreases sharply when the vector **q** exceeds the characteristic dimension of the region of the "slow" phonons in the corresponding direction. We have thus returned to Akhiezer's concept of a constant phonon viscosity, but only for the "slow" phonons, which freeze out rapidly with decreasing temperature, as is reflected in the temperature dependence of $C_S(T)$.

Breaking up the integral (2.2) into two terms corresponding to the regions $ql < 1$ and $ql > 1$, we can easily obtain, taking (2.23), (2.26), and (2.27) into account, an estimate of the coefficient of the dragging of the dislocations[3)]:

$$B \approx \frac{(\alpha b)^2}{64\pi}\left\{\frac{CT}{\chi}\ln\left(\frac{c}{v}\right)^s + \frac{1}{\pi}\hbar k_D^5\left[f_1\left(\frac{T}{\Theta}\right) + \lambda f_2\left(\frac{T}{\Theta_s}\right)\right]\right\}; \qquad (2.30)$$

here

$$\lambda \sim \frac{10g}{\beta^2} k_D l, \quad f_2(x) = \frac{1}{x}\frac{e^{1/x}}{(e^{1/x}-1)^2}, \qquad (2.31)$$

Θ_S is the characteristic temperature of the "slow phonons" ($\Theta_S \sim \Theta$), s is a numerical coefficient on the order of unity (s = 0 for a screw dislocation and s differs from zero in the case of an edge dislocation only at not too low temperatures, when the Umklapp processes are not small[26]). The first term of (2.30), corresponds to thermoelastic damping and coincides with the Lothe estimate[4] (2.11) at $\alpha = \alpha_B = \gamma$. The second term is an estimate of the contribution made to the slowing down by the scattering processes. At $\alpha = \alpha_A = M/4G$ this term coincides with the phonon-wind estimate (2.19) and (2.24), previously obtained under the assumption of high dislocation velocities. The third term in (2.30) corresponds to relaxation of the "slow" phonons.

Expression (2.30) enables us to compare the relative contributions of the dissipative processes of the types considered above to the dragging of the dislocation. At low temperatures, when the first and third terms are exponentially small, scattering processes predominate. At high temperatures, when $f_1(T/\Theta)$, $f_2(T/\Theta)$, and CT are practically linear in temperature, and the mean free path of the phonons varies in inverse proportion to the temperature[4)], $l = l_\Theta(\Theta/T)$, expression (2.30) has a structure (with conservation of the order of the terms)

$$B = B_0\left[A\left(\frac{T}{\Theta}\right)^2 + \frac{T}{\Theta} + 2\beta\lambda_\Theta\right]; \qquad (2.32)$$

Here $A \sim 10^{-1}(r_0/l_\Theta)\ln(c/v)$, and λ_Θ is the value of λ at T = Θ. Since $A \ll 1$ ($A \sim 10^{-1}$ at $l_\Theta \approx 5r_0$ and $c \approx 10^2 v$), the thermoelastic processes are masked by the phonon wind but can be separated in principle at high temperatures by means of the quadratic temperature dependence. To this end it is necessary, however, to increase appreciably the experimental accuracy. The quantity $2\beta\lambda_\Theta$ turns out to be of the order of unity, i.e., the relaxation of the "slow" phonons makes a contribution comparable with the contribution of the phonon wind to the dragging of the dislocations, if the region of the "slow" phonons occupies in the first Brillouin zone a relative volume $g \sim 10^{-1} r_0/l_\Theta$, which is perfectly realistic. Generally speaking, in view of the complexity of the true phonon spectra, the quantity λ_Θ is rather difficult to calculate. In any case, it would be difficult to get along without a computer. Instead, we shall regard λ_Θ, as we do the Debye temperature, as a phenomenological parameter to be determined from experiment. We shall show in Chap. 4 that the function

$$f_1\left(\frac{T}{\Theta}\right) + \lambda_\Theta\frac{\Theta}{T} f_2\left(\frac{T}{\Theta_s}\right)$$

describes adequately the experimental B(T) curves at reasonable values of λ_Θ.

f) Contribution to the dissipation of the optical phonons. So far, when analyzing the dynamic dragging of the dislocations, we have considered only the acoustic branches of the phonon spectrum, which are present in all the crystals. Yet crystals that have more than one atom per unit cell contain in the phonon spectrum also optical modes, for which $\omega(0) \neq 0$. A typical dispersion curve for optical phonons is shown schematically in Fig. 6.

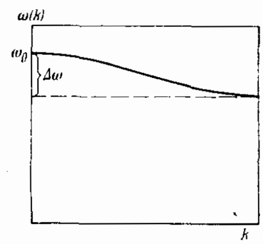


FIG. 6. Typical form of the dispersion curve ω(k) for optical phonons.

The contribution of the optical phonons to the damping of the long-wave part of the packet (2.1) ($ql < 1$) is limited by the thermoelastic losses, and its estimate is contained in the first term of (2.30), since it is implied that the specific heat C is determined by all the branches of the phonon spectrum. When analyzing the short-wave region $ql > 1$ by means of formula (2.27) it is convenient, as before, to distinguish between the "fast" and "slow" phonons, which interact in a qualitatively different manner with the moving dislocation. The interaction of the "fast" phonons with the dislocation is of the scattering type, and the contribution of these phonons to the dissipation can be estimated from formula (2.18). The "slow" phonons relax like a gas with effective viscosity (2.29) multiplied by $(\Theta/\Theta_0)^4$, where $\Theta_0 = \hbar\omega_0/k_B$.

Inasmuch as low temperatures ($T \ll \Theta_0$) the density of the optical phonons is exponentially small, a noticeable effect can be expected only at sufficiently high temperatures. Not being interested in temperatures at which the effect is certainly known to be small, it is convenient to confine oneself only to the temperature region $T \gg \hbar\Delta\omega/k_B$; this simplifies the analysis appreciably. It can be assumed here that $vk_D \ll \Delta\omega \ll \omega_0$.

As expected, calculation predicts an exponential smallness of the effect at low temperatures $T \ll \Theta_0$, but at high temperatures the dependence is of the form (2.32). The contribution of the "fast" phonons amounts to a fraction on the order of $10^{-1}(\Theta/\Theta_0)^4\omega_D/\Delta\omega$ of the phonon wind, while the contribution of the "slow" phonons differs from the corresponding term in (2.30) by a factor $10^{-1}(\Theta/\Theta_0)^4$. Usually Θ_0 noticeably exceeds the Debye temperature, and the dissipative processes in the system of acoustic phonons should prevail over the damping in the optical modes of the spectrum. A possible exception is crystals containing "soft" optical modes in the phonon spectrum. In particular, under conditions of a phase transition of the displacement type, "soft" modes usually appear for which $\Theta_0 \to 0$ as $T \to T_c$. Accordingly, the optical component of the viscous slowing down of the dislocations should increase without limit near the transition temperature. The strong dependence of the dragging coefficient B on the quantity Θ_0 ($B \propto \Theta_0^{-4}$) gives grounds for hoping that the predicted effect will be readily observed in experiment. The resonant decrease of the dislocation mobility near the transition point has a paradoxical character against the background of the general softening of the crystal (near T_c, the crystal can behave as a liquid with respect to certain oscillation modes).

3. PHONON DRAGGING MECHANISMS DUE TO EXCITATION OF NATURAL DEGREES OF FREEDOM OF MOVING DISLOCATIONS

In the preceding chapter we have considered the phonon dragging of a linear dislocation that moves uniformly as a unit. Yet allowance for the internal degrees of freedom of the dislocation leads to a number of qualitatively new effects. Thus, perturbations of the linearity of the shape and of the uniformity of the dislocation motion in the thermal field of the lattice give rise to induced radiation of phonons by the dislocation, and this radiation increases the dissipation (the flutter effect). The motion of a dislocation in the periodic potential due to the discrete character of the crystal leads to periodic changes in the structure of the core and to oscillations of the velocity of the dislocation. The corresponding changes of the elastic field and of the dislocation energy give rise to emission of waves (radiative dragging) and to Raman scattering of the phonons by the oscillations of the elastic field. Additional phonon radiation is produced also when a dislocation moves near various types of lattice defects, which produce local distortion fields. Under certain conditions, the foregoing effects turn out to be appreciable against the background of the dissipative processes investigated above.

a) Flutter effect. So far we have dealt only with those phonon-dislocation interactions which are due to nonlinear properties (anharmonicity) of the crystal. However, the nonlinear properties of the medium are not the only cause of interaction between a dislocation and the phonon subsystem. Even in a harmonic crystal, phonons can be scattered by dislocation as a result of the so-called flutter effect. The dislocation has its own degrees of freedom, which are on a par with the remaining degrees of freedom of the crystal. Vibrating in the thermal motion of the lattice, a dislocation, as any other source of internal stresses, radiates elastic waves. In other words, in a crystal with a dislocation the wave function of the phonon should be a superposition of a plane wave and a wave diverging from the dislocation, and this can be described in terms of phonon scattering.

This effect can be estimated by considering the induced oscillations (flutter) of a dislocation in the alternating stress field connected with the individual phonons. Within the framework of this approach, Nabarro[16] calculated the total cross section of the flutter effect for long-wave phonons that are normally incident on a dislocation at rest, but reached the erroneous conclusion that the total momentum of the radiation induced as the dislocation moves is zero. Lothe has shown[4,17] that the presence of aberrational symmetry of the induced radiation leads to an effective dragging of the dislocation, proportional to its velocity. A more detailed investigation of the flutter effect for a dislocation at rest was carried out by Ninomiya[29] as applied to the problem of the influence of dislocations on the thermal conductivity. The differential cross section of the flutter effect calculated in[29] was used by Al'shitz and Sandler[30] and by Ninomiya[31] for a consistent calculation of the contribution of this mechanism to the dragging of the dislocation and to assess its relative role[30] against the background of the phonon wind.

The flux density of phonons in a state α, incident on a dislocation, is equal to cn_α. Accordingly, the number of phonons scattered per unit time from the state α into the state β is $cn_\alpha\sigma_{\alpha\beta}$, where $\sigma_{\alpha\beta}$ is the differential cross section of the flutter effect on the moving dislocation. The quantity $\sigma_{\alpha\beta}$ can be obtained from the corresponding expression for $\sigma^0_{\alpha\beta}$ of the dislocation at rest[29] by means of the usual Doppler frequency shift: $\omega'_\alpha = \omega_\alpha + \mathbf{v}\cdot\mathbf{k}$. The requirement contained in $\sigma^0_{\alpha\beta}$, that the quantity ω'_α be conserved during the scattering process, means that in in each scattering act the energy transfer is $\hbar(\omega_\alpha - \omega_\beta) = \hbar\Omega_\mathbf{q}$. The energy dissipation per unit time is then determined by the expression

$$D = \sum_{\alpha,\beta} \hbar\Omega_\mathbf{q} cn_\alpha\sigma_{\alpha\beta}. \quad (3.1)$$

Taking the foregoing into account, calculation by means of formula (3.1) entails no difficulty. The dissipation (3.1) corresponds to a dragging coefficient[30]

$$B = \frac{\hbar k_D^3}{2\pi^2} f_3\left(\frac{T}{\Theta}\right). \quad (3.2)$$

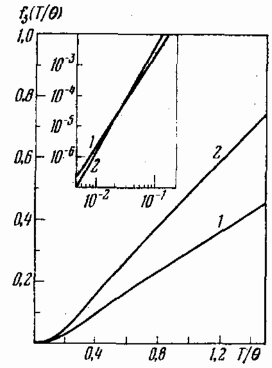


FIG. 7

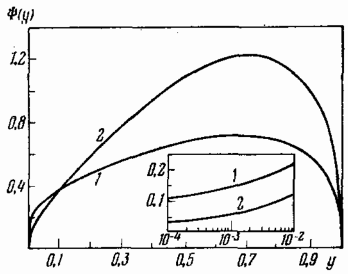


FIG. 8

FIG. 7. Plot of the function $f_3(T/\Theta)$, which determines the temperature dependence of the flutter effect for a screw dislocation (1) and an edge dislocation (2).

FIG. 8. Plot of the function $\Psi(y)$ for a screw dislocation (1) and an edge dislocation (2).

The temperature dependence of the effect is determined by the function (Fig. 7)

$$f_3(x) = x^3 \int_0^{1/x} \frac{dt\, e^t t^3}{(e^t-1)^2} \Psi(xt); \tag{3.3}$$

Here $\Psi(y)$ is a logarithmically slow function[5)] shown graphically in Fig. 8. In the case of an edge dislocation we have

$$\Psi(y) \approx \frac{1,2}{1+\pi^{-2}\ln^2(y^{-2}-1)}. \tag{3.4}$$

It should be borne in mind that formula (3.2) is exact in the sense that the expression for the phonon scattering cross section, which was used in its derivation, was obtained without perturbation theory. Nonetheless, this estimate contains another source of errors, since no account was taken in it of the finite dimensions of the dislocation core. Allowance for these dimensions, which is of no importance at low temperatures when the average phonon wavelength greatly exceeds the dimensions of the core, can greatly decrease the amplitude of the effected high temperautres, as we have seen with the phonon wind as an example. Ninomiya [31], while noting that the dislocation core lowers the estimate of B at room temperature by approximately one order of magnitude, assumes, however, that this decrease can be offset by taking additional account of the resonant scattering of the phonons due to the excitation of local dislocation vibrational modes. It appears that Ninomiya [31] nevertheless overestimates the role of the logarithmic singularities due to the resonant scattering. At any rate, the estimate of B by means of formula (3.2), obtained in [30] with allowance for resonance effects, differs by a factor on the order of unity from Ninomiya's estimate in which these effects are neglected. Thus, one cannot speak seriously of any compensation, and the comparison made in [31] between the results of the calculations of B by a formula of the type (3.2), which does not take into account the existence of the dislocation core, with values of B measured at room temperature, is hardly meaningful. One can confine oneself to the "coreless" approximation only to the extent to which, as we shall show, the flutter is noticeable against the background of other mechanisms of dislocation dragging, only at low temperatures when this approximation is quite satisfactory.

It follows from (3.2) and (3.3) that at high temperatures the flutter effect, like the phonon wind, varies linearly with temperature: $B \propto T$. At low temperatures, the linear dependence gives way to a cubic dependence $B \propto T^3$, corresponding to the slowest of all the considered mechanisms of the decrease of the dragging with decreasing temperature. It is seen from Fig. 7 that, just as in the case of phonon wind, the dragging becomes linear quite rapidly—in fact, even at temperatures much lower than the Debye temperature.

The estimate (3.2) enables us to compare the flutter effect with phonon wind. At low and high temperatures, the ratio of the dragging coefficients B_{fl}/B_W should tend asymptotically to the values

$$\frac{B_{fl}}{B_W} = \left(\frac{T_0}{T}\right)^2 \quad (T \ll \Theta), \qquad \frac{B_{fl}}{B_W} = d_0 \approx \text{const} \quad (T \gg \Theta), \tag{3.5}$$

where the quantities T_0 and d_0 are of the order of[6)]

$$T_0 \sim \frac{10G}{M}\frac{\Theta}{k_D b}, \quad d_0 \sim \left(\frac{10G}{M}\right)^2. \tag{3.6}$$

Estimating the ratio M/G from formula (2.21), we can show that for copper, for example, $T_0 \approx 15°K$ and $d_0 \approx 0.1$.

Thus, from among the mechanisms considered so far, the flutter effect should predominate at low temperatures, and a combination of phonon wind and relaxation of the slow phonons should predominate at high temperatures. Incidentally, one can conceive of weakly-anharmonic crystals ($d_0 > 1$) in which the flutter effect predominates over the phonon wind also at high temperatures.

A comparison of the analogous mechanisms [27,32,33] of the dragging of kinks leads to qualitatively similar conclusions [25,27].

b) Radiation friction. Owing to the discrete character of the lattice, when a dislocation moves in a crystal its atomic configuration and elastic energy experience periodic changes. The corresponding relief is called in the literature the Peierls relief. The quasistatic properties of this relief were investigated in a number of studies using the Peierls-Nabarro model [34-36] and also the Frenkel-Kontorova model [37-40]. To set into motion a dislocation that lies in the valley of this relief, it is possible to apply to it a certain "starting" stress, called the Peierls stress σ_P. In the case of stationary motion of the dislocation, however, the influence of the relief does not reduce by far to the static resistance of the lattice. In fact, in the absence of dynamic energy losses the motion of the dislocation over the Peierls relief would be ensured by successive conversion of the potential energy into kinetic energy and vice versa, and in order to maintain this motion there would be no need for an external force at all. The role of the relief in this sense reduces to stimulation of dynamic losses, since periodic changes of the configurations of the dislocation core, and the non-uniformity of its motion over the relief, should lead to radiation of the elastic waves by the dislocation, i.e., to radiation friction. This dissipation mechanism, due exclusively to the discrete character of the lattice, is preserved also at the lowest temperatures, in contrast to the phonon effects considered above, which freeze out with decreasing temperature. In the literature, the contributions of the configuration oscillations ("breathing" of the core) and dynamic oscillations (variations of the velocity) to the radiation friction were investigated independently, as individual dissipation channels. This separation can be justified in principle, although it seems aribtrary at first glance. We formulate below the condition under which this separation is meaningful.

The radiation of elastic waves by a dislocation, due to the periodic changes of the form of the core of the dislocation, was considered in the simplest formulation back by Frenkel and Kontorova [37]. Subsequently, the model of Frenkel and Kontorova was used by Weiner [41] and by Atkinson and Cabrera [42] for a more general analysis of the problem of radiative dragging of a uniformly moving dislocation. According to [41], radiation friction, up to near-sonic velocities, is practically independent of the velocity and can be described by the so-called dynamic Peierls stress σ_{DP} which amounts to only a small fraction of the static stress σ_P. Only at velocities close to the velocity of sound ($v \gtrsim c/2$) does the radiative dragging of the dislocation increase abruptly and approach values exceeding the stress σ_P.

Atkinson and Cabrera [42], analyzing the same problem in a somewhat more accurate formulation, obtained for high dislocation velocities ($v > c/3$) results analogous to those of [41]. However, according to [42], with decreasing velocity the dislocation excites an ever increasing number of phonon modes. In this velocity region ($v < c/3$), the authors observed a series of infinite resonances in the radiation, and the series condensed towards lower velocities. The resonant situation is realized whenever the dislocation velocity coincides with the group velocity in one of the elastic waves excited by it in the lattice. Incidentally, as noted in [42], in the region of the resonance points the amplitude of the oscillations of the atoms of the chain should increase without limit, in contradiction to the initial assumption of the theory. This seems to indicate that there are no stationarity solutions near the resonances, and possibly also in the entire region of low velocities. Unfortunately, the region where the stationary solutions exist was not determined in [42]. The authors of [42] are inclined themselves to assign a physical meaning only to the results pertaining to the high-velocity region $c/3 \lesssim v < c$. They attribute the singularities of the radiative dragging in the region of "low" velocities ($v < c/3$) to the one-dimensional character of the Frenkel-Kontorova model, and to the use of a piecewise-harmonic potential for the atoms of the chain.

Notice should also be taken of the work of Kosevich and Margvelashvili [43], who investigated acoustic and electromagnetic emission on the uniform motion of a dislocation in an ion crystal under the assumption that the radiation occurs at the fundamental frequency $\Omega = 2\pi v/a$. However, it seems that in problems of this type, the neglect of radiation at high harmonics $\Omega_n = \Omega n$ can in general not be justified.

The problem of the onset of radiation-friction resonances when perturbation sources move in periodic structures was investigated in rather general form (albeit still in the same approximation of a piecewise-harmonic atomic potential) by Rogula [44], who traced, in particular, the influence of the dimensionality of the problem on the character of the singularity at resonance. According to [44] the singularities at the resonance points become smoothed out with increasing dimensionality of the problem. Thus, in the two-dimensional case, an example of which is an infinite straight-line dislocation, we are dealing not with a square-root singularity of the type $(v - v_p)^{-1/2}$, which is typical of the one-dimensional problem [42], but with the discontinuities and logarithmic singularities $\propto \ln|v - v_p|$. In the three-dimensional case, the radiative dragging as a continuous function of the velocity and only the derivative of the dissipative function has discontinuities. It should be noted, however, that the question of the conditions for the existence of stationary solutions was not investigated in that study, although the author did show that stationary motion at resonant velocities is impossible. One cannot exclude, in principle, the possibility that the resonant radiation has no physical meaning at all, since it pertains to the velocity region where the stationary motion is not realized, and consequently it is impossible to formulate a problem in which resonances arise.

It appears that Rogula's work remained unknown to a number of researchers engaged in similar problems. A few years later [44], two papers of similar content were published by Celli and Flytzanis [45] and by Ishioka [46], in which, without mentioning Rogula's results, the two-dimensional problem was again considered as applied to radiative dragging of an infinite linear screw dislocation. The authors of [45,46] generalized, to include the case of uniform motion of the dislocation, the model of Maradudin [47a], who proposed a method of discretely describing an immobile screw dislocation in a cubic lattice (with a piecewise-harmonic atomic potential). This model constitutes in essence a two-dimensional modification of the Frenkel-Kontorova model. Naturally, the principal results obtained in [45,46] are related to a considerable degree to the conclusions of [44] and do not differ qualitatively from the results of [42], and therefore will not be reported here. We note only that these papers also failed to explain the physical meaning of the resonant radiation, since the question of the region where stationary solutions exists was not investigated.

A separate problem of undoubted physical interest is the possibility of near- and supersonic motion of dislocations. It is known, for example, that in the continual model, the dislocation energy becomes infinite when it passes through the sonic barrier. However, neither in [42] nor in [44-46] were any singularities whatever observed in the radiative dragging of the dislocations in the vicinity of the speed of sound. Earmme and Weiner [47b], who recently turned to this problem, revised the results of Atkinson and Cabrera for high-velocity dislocations. Unlike their predecessors, Earmme and Weiner [47b] investigated the region of existence of stationary solutions and have shown that the stationarity conditions are violated even in the subsonic region, when the dislocation velocity exceeds a certain critical value $v_B < c$. Thus, the continual theory withstood one more test of the correctness of the qualitative predictions, although the problem undoubtedly still calls for further study.

An interesting formulation of the problem of radiative slowing down of dislocations was proposed by Flytzanis and Celli [48]. Within the framework of the approach developed in [45,46], they attempted to consider the radiation friction and the flutter effect from a unified point of view, i.e., to take additional account of the non-uniform motion of the dislocation in the thermal field of the lattice. Unfortunately, the complexity of the problem forced them to resort to major simplifications, which yielded an approximate solution only for near-sonic and supersonic dislocations. Incidentally, following the work of Earmme and Weiner, even this result cannot cause any doubts. Nonetheless, it appears that one can agree with the general qualitative conclusion of [48] that the radiation friction can be noticeable against the background of the phonon dragging only for near-sonic dislocations at low temperatures.

In a recent paper[49], Ishioka returned to the one-dimensional dislocation model, including in consideration also the damping of the natural oscillations in the system and the anharmonic character of the atomic potential in the chain. One of the most interesting results of the work was the proof that the hitherto employed piecewise-harmonic approximation of the atomic relief is not suitable for the description of the motion of the dislocation at lower velocities, or else it leads to the absence of stationary solutions of the problem in this velocity region. Ishioka has shown that this difficulty does not arise in the case of smooth potentials having a continuous first derivative. The problem was solved numerically in[49] for a sinusoidal potential and for two types of smoothly-joined piecewise-parabolic potentials. The radiation friction turned out to be a smooth function that decreases monotonically with decreasing dislocation velocity. This seemed to solve the question of the existence (more accurately, of the absence) of resonant radiation at low dislocation velocities, the investigation of which has been the subject of so many studies. In essence, we are dealing with a rather typical situation, stemming from the fact that frequently the dissipative function corresponding to the stationary solution is much easier to find than the solution itself. Accordingly, whenever the dissipation is determined without investigating the region of the existence of the stationary solutions, there is a danger of obtaining all kinds of fictitious effects pertaining to the region where there are no stationary solutions.

According to[49], at near-sonic velocities the dislocation field has a clearly pronounced dynamic wave zone, and the dislocation radiates intensively elastic waves during the course of its motion. The radiation slowing down of the dislocation then greatly exceeds the static Peierls stress. When the dislocation velocity decreases, the "relativistic" effects increase and the level of the radiation friction drops abruptly. Accordingly, the dislocation field becomes closer and closer to the quasi-static field that moves uniformly with the dislocation. An important circumstance is that the region of the quasi-static motion of the dislocation begins at sufficiently high velocities, corresponding to a dislocation kinetic-energy level much higher than the Peierls level, when the perturbation of the uniformity of the dislocation motion can still be regarded as negligible. It is this which makes it possible to separate the contribution of the configuration field oscillations of the dislocation from the contribution of the dynamic oscillations to the radiation friction, since the former are significant only at high velocities, when the dislocation motion is practically uniform, whereas the latter come into play in the dissipation at low velocities, when the first dissipation channel can be regarded as suppressed.

For a correct solution of the problem of the radiation friction due to the non-uniform motion of the dislocation, it is necessary to determine in self-consistent fashion the law of dislocation motion in a periodic potential field, with allowance for the reaction of the radiation. This approach was first employed in a paper by Al'shitz[50], who made, however, the traditional error referred to above, namely he calculated the radiative dragging of the dislocation, but did not investigate the conditions for the existence of stationary solutions of the problem. This error was corrected subsequently in a paper by Al'shitz. Indenbom and Shtol'berg[51]. In[50,51] we investigated in the continual approximation the stationary motion of a screw dislocation over a Peierls relief under the influence of a constant external force f that compensates for the radiation losses. The exact solution of the problem was obtained for a piecewise-parabolic Peierls relief[50,51] and an approximate solution was obtained for a sinusoidal relief[51]. General expressions were obtained for the law of dislocation motion and for the radiation stress $\sigma = f/b$. An analysis of these expressions shows that at high velocities v, when the kinetic energy of the dislocation greatly exceeds the Peierls energy, the presence of the relief disturbs the uniform motion of the dislocation only insignificantly, and the radiation is mainly at the first harmonic, while the radiation friction, in accordance with an estimate by Hart[52], decreases in proportion to v^{-2}. With decreasing dislocation velocity, the degree of nonuniformity of its motion increases, and accordingly the radiation losses increase, and the radiation at a higher harmonics becomes more and more effective. A decrease of the average velocity is possible only to a certain critical value $v_c \sim c\sqrt{\sigma_P/G}$, and the minimum possible average dislocation velocity corresponds to motion in which the dislocation at the crest of the relief has zero kinetic energy.

A critical velocity exists also under conditions of viscous dissipation, when the dislocation is acted upon additionally by a certain viscous dragging force $f_f = -B\dot{x}$, but for not too high values of the viscosity B. The critical velocity decreases with increasing viscosity and, starting with a certain B_c, it vanishes—the stationary motion at $B > B_c$ is realized at all velocities v. Without allowance for the viscosity, the dynamic dragging is determined only by the radiation friction and is characterized by a decreasing function of the velocity, corresponding to instability of the stationary motion. The viscous dissipation stabilizes the motion by adding to the dragging force a term linear in the velocity. At $B > B_c$, an effect of the "dry friction" type should occur, namely, the stress σ approaches the static Peierls stress σ_P with decreasing velocity v, and does not vanish when the velocity tends to zero. The asymptotic form of the function $\sigma(v)$ at $B \gg B_c$ is described by the simple formula

$$\sigma(v) = \sigma_P \operatorname{cth} \frac{b\sigma_P}{Bv}, \tag{3.7}$$

which illustrates clearly the phenomenon of "dry friction."

A similar problem was solved[53] for tangential motion of a kink along a dislocation with allowance for a secondary Peierls relief. It was shown that all the qualitative regularities noted above hold true also in the case of a kink.

c) Raman scattering of phonons. As a result of periodic changes in the configuration of the core and the velocity of the dislocation, configurational and dynamic oscillations of the elastic field of the dislocation, at the fundamental frequency $\Omega = 2\pi v/a$ and at its overtones $\Omega_n = \Omega n$, are produced respectively as the dislocation moves along the Peierls relief. The phonons are scattered by this field in an inelastic (Raman) manner, experiencing an energy change $\Delta E = \pm\hbar\Omega_n$. The predominance of the Stokes component of the scattering ($\Delta E > 0$) over the anti-Stokes component ($\Delta E < 0$) determines the dissipation of the energy and the effective dragging of the dislocation. This mechanism was first considered by Al'shitz[50,54].

If $W^{\pm}_{\alpha\beta}$ is the number of transitions, per unit time, of the phonons from the state α to the state β with absorption (+) or emission (−) of an energy $\hbar\Omega_n$, then the power of the dissipative losses is determined by the formula

$$D=\frac{1}{2}\sum_n\sum_{\alpha,\beta}\hbar\Omega_n(W^+_{\alpha\beta}-W^-_{\beta\alpha}). \qquad (3.8)$$

The values of $W^{\pm}_{\alpha\beta}$ are given by the usual quantum-mechanical expression

$$W^{\pm}_{\alpha\beta}=\frac{2\pi}{\hbar^2}|V_{\alpha\beta}|^2 n_\alpha(n_\beta+1)\,\delta(\omega_\alpha-\omega_\beta\mp\Omega_n), \qquad (3.9)$$

where $V_{\alpha\beta}$ is the matrix element of the transition. Thus, we obtain from (3.8) and (3.9) a formula of the type (2.18)

$$D=\frac{\pi}{\hbar}\sum_n\sum_{\alpha,\beta}\Omega_n^2|V_{\alpha\beta}|^2\frac{\partial n_\alpha}{\partial\omega_\alpha}\delta(\omega_\alpha-\omega_\beta). \qquad (3.10)$$

It can be shown that the quantity $V_{\alpha\beta}$ has a structure of the form

$$V_{\alpha\beta}=h_n b\left(\frac{\mathbf{qv}}{v}\right)\Gamma_{\alpha\beta}, \qquad (3.11)$$

where the coefficient h_n is a dimensionless parameter proportional to the amplitude of the oscillations of the elastic field of the dislocation. Taking (3.11) into account, we can easily verify that expression (3.10) differs from (2.18) by a constant temperature-independent factor. The values of this factor I, which characterizes the relative role of the Raman scattering in comparison with the phonon wind, turn out to be different for the configuration and for the dynamic oscillations:

$$I_{\text{con}}\approx 5\pi^3\left(\frac{r_0}{b}\right)^2\frac{\sigma_P}{G},\quad I_{\text{dyn}}\approx\left[\frac{\sigma_P}{G}\frac{c^2}{v^2}\left(\ln\frac{c}{v}\right)^{-1}\right]^2. \qquad (3.12)$$

It follows from (3.12) that the Raman scattering of phonons by configuration oscillations of the dislocation field introduces into the energy dissipation a contribution that is noticeable against the background of the phonon wind only in crystals with a sufficiently high Peierls relief, namely, at $r_0 = 3b$ we have $I_{\text{con}} \gtrsim 1$ if $\sigma_P/G \gtrsim 7\times10^{-4}$. The fact that the numerical estimates based on formula (3.10) were found in [50,54] to agree with the experimental data for certain crystals at $\sigma_P/G = 5\times10^{-5}$ is due to the use in [50,54] of the usual formulas of the continual theory of dislocations, in which it is assumed that $r_0 = 0$. Using the transition from formula (2.20) to formula (2.24) as an example, we have seen that allowance for finite character of the dislocation core leads to a decrease in the amplitude of the effect by a factor $k_D r_0$, i.e., by approximately one order of magnitude.

The relative role of the dynamic oscillations can be easily established if it is noted that I_{dyn} is of the order of the square of the ratio of the Peierls energy to the kinetic energy of the dislocation. In the dynamic velocity region, which we are considering, I_{dyn} is therefore small, thus indicating that this dissipation channel is negligible in comparison with the phonon wind. However, recognizing that the dynamic oscillations lead to a dragging that is essentially nonlinear in the velocity ($I_{\text{dyn}} \propto v^{-4}$), we can hope to separate their contribution, using this nonlinearity, against the background of the viscous slowing down.

d) Influence of impurity. When the dislocations overcome the local fields of the impurity centers and other lattice defects, the uniformity of the dislocation motion becomes disturbed and phonons are radiated, so that the level of the radiation friction increases. In particular, Ookawa and Yazu [55] investigated the mechanism of the emission of elastic waves in the field of impurity centers, a mechanism similar to the mechanism of radiation friction for dynamic oscillations of the field of a dislocation moving along a Peierls relief. According to their calculation, the dragging force is relatively small, decreases with velocity like v^{-1} at high dislocation velocities, and vanishes as v tends to zero. Incidentally, the estimate obtained for slow dislocations without a self-consistent determination of the law of dislocation motion cannot be regarded as reliable.

Another mechanism of dynamic dragging of dislocations by impurity centers can be connected with excitation of local or quasilocal vibrations of impurity atoms [56-59], an effect analogous to the well known Bohr losses in electrodynamics. In the most general formulation, this effect was considered by Kosevich and Natsik [58,59] within the framework of the phenomenological theory previously developed by them [9], in which the dragging of the dislocations is connected with the dispersion of the elastic moduli. The authors have solved the problem of scattering of the dislocation wave packet (2.1) by atoms of a heavy interstitial impurity, taking into account the possibility of exciting quasilocal impurity oscillations. According to [59], the dragging force has a sharply pronounced maximum at a dislocation velocity on the order of $a\omega_L$, where ω_L is the frequency of the quasilocal level. At low velocities, the dragging force increases with velocity in proportion to v^3, and at high velocities it decreases like v^{-1}. If the impurity does not excite any quasilocal vibrations, then the dragging force is proportional to v^3 in the entire velocity interval. The entire effect is proportional to the square of the ratio of the impurity-atom mass to the matrix-atom mass.

The influence of the impurity on the dynamic mobility of dislocations can be due not only to the phonon mechanisms, but also to dissipative processes connected with the diffusion mobility of the impurity in the field of a moving dislocation. This problem was investigated by many workers [59-68] from different points of view and in various approximations. For example, the diffusion dragging of fast dislocations moving with practically no atmosphere was considered in the same paper of Kosevich and Natsik [59]. Notice should also be taken of a recent paper by Lyubov and Altundzhi [58], who investigated the dragging of slow dislocations moving together with their impurity atmospheres. As a rule, the contribution of the diffusion mechanisms to the dragging of the dislocations is negligible against the background of the phonon effects.

The influence of the structural imperfections of the crystal on the dynamic energy losses can also be due to the buildup of oscillations of segments of pinned dislocations (so-called dislocation "scaffold") in the elastic field of a moving dislocation (see, for example, the review of Indenbom and Orlov [69]). The resultant additional dissipation by the vibrating segments is limited with the same mechanisms as the direct dragging of the dislocation. Therefore the corresponding increment to the drag coefficient should have the same temperature dependence as the main term, and its role reduces to a renormalization of the absolute value of the effect. Natsik and Minenko [70] carried out a quantitative estimate of this effect and have shown that the renormalization can be noticeable at reasonable densities of the dislocation "scaffold."

4. EXPERIMENTAL DATA ON THE DYNAMIC DRAGGING OF DISLOCATIONS

To measure the dynamic dragging of dislocations, one can use various methods based on the analysis of the above-barrier motion of fast dislocations, on the determination of the damping of the dislocation segments that oscillate between the pinning centers, and an estimate of the macroscopic viscosity of the crystals that undergo high-velocity plastic deformation.

a) Measurement of the mobility of individual dislocations. A small number of fresh dislocations is introduced (usually by local deformation of the sample) in a crystal with low dislocation density and the path Δl of the dislocations under the influence of a stress pulse of known amplitude σ and duration Δt is observed. The initial and final positions of the dislocations are usually obtained by selective etching of the surface of the crystal before and after application of the pulse. As a rule, the duration of the pulse exceeds by several orders of magnitude the characteristic time $\tau_0 \sim Gb^2/Bc^2$ of establishment of the stationary motion of the dislocation, so that the measured values of Δl turn out to be proportional to Δt, and their ratio $v = \Delta l/\Delta t$ is a measure of the average velocity of the dislocation at a given level σ of the external stress. This method makes it possible to plot the mobilities of individual dislocations for different crystals in a wide range of velocities (see Figs. 1 and 2). At high velocities the function $v(\sigma)$ is linear, and its slope characterizes the level of the dynamic dragging of the dislocations and makes it possible to determine the value of the dragging coefficient B:

$$B = \lim_{\sigma \to \infty} \frac{b\sigma}{v(\sigma)}. \tag{4.1}$$

In practice, the measurement accuracy is limited by the considerable scatter (on the order of 30%) of the path lengths of the dislocations, owing to the presence of internal stresses and various structural imperfections in the crystal. Nonetheless, a direct determination of the mobility of the dislocations is the most reliable method of investigating the dynamic dragging, and it is precisely this method that yielded the bulk of the hitherto accumulated experimental material [71-86].

b) Amplitude-independent internal friction. When high-frequency ultrasound (10^6—10^8 Hz) is passed through a crystal, an amplitude-independent internal friction is observed. The frequency dependence of the damping has a broad maximum that shifts towards lower frequencies with increasing temperature. In this region, the intensity of the internal friction is determined mainly by the viscous losses that are produced when the dislocation segments vibrate between the pinning points. According to the theory of Granato and Lucke [87], in which the dislocation is regarded as a string with a certain constant linear tension $T \sim Gb^2$, which in turn is determined by the linear energy of the dislocation, the dependence of the amplitude-independent internal friction on the frequency is described by the expression

$$Q^{-1} = k\Delta_0 \rho_d L \frac{\omega/\omega_0}{1+(\omega/\omega_0)^2}; \tag{4.2}$$

where k is an orientational factor, ρ_d is the dislocation density, L is the effective length of the segment, $\Delta_0 = 8Gb^2/\pi^4 T$, and $\omega_0 = \pi^2 T/BL^2$. The last relation enables us in principle to determine B from the position of the maximum $\omega = \omega_0$. Since, however, ω_0 depends on the length L and is therefore sensitive to the method used to average the segments over the lengths, the dragging coefficient B is usually determined from the asymptotic form of the descending part of the $Q^{-1}(\omega)$ curve, which does not depend on L:

$$\lim_{\omega \to \infty} (\omega Q^{-1}) = \frac{8}{\pi^2} kGb^2 \frac{\rho_d}{B}. \tag{4.3}$$

The dragging coefficient B was measured by this method in a large number of crystals [19,88-104]. The accuracy of the method is limited by the errors in the determination of the dislocation density ρ_d, which are usually appreciable.

c) High-velocity deformation. Another indirect method of determining the dragging coefficient of the dislocations is connected with macroscopic experiments on the deformation of crystals at high velocities. The rate of plastic deformation $\dot{\epsilon}$ can be expressed in terms of the density ρ_m of the mobile dislocations and in terms of their average velocity v:

$$\dot{\varepsilon} = b\rho_m v. \tag{4.4}$$

To overcome the forces of the dynamic dragging of the dislocations it is necessary to apply a stress

$$\sigma_B = \frac{Bv}{b} = B\frac{\dot{\varepsilon}}{\rho_m b^2}. \tag{4.5}$$

In some cases the contribution of σ_B to the total resistance of the crystal to plastic deformation can be estimated from the dependence of the flow stress on the deformation rate (in the general case this dependence can be due also to the influence of the stress on the rate at which the barriers are overcome). In particular, in the case of deformation of crystals with high velocity ($\dot{\epsilon} \gtrsim 10^3$ sec^{-1}), when it can be assumed that all the barriers are overcome without participation of thermal fluctuations, the experimentally observed linear dependence of the yield point on the deformation rate

$$\sigma_y = \sigma_0 + \alpha_0 \dot{\varepsilon} \tag{4.6}$$

can be interpreted as a direct manifestation of the dynamic dragging $\sigma_B = \alpha_0\dot{\epsilon}$, and one can estimate B with the aid of (4.5) and (4.6) from the value of α_0 and from the density ρ_m of the mobile dislocations [105-110]. This method is apparently not very reliable, since only the order of magnitude of ρ_m is usually known.

Analogous estimates of B can be obtained by reducing the data on the energy dissipation in shock waves that imitate deformation by explosion [111]. In this case, however, considerable difficulties arise with the determination of the density of the mobile dislocations, since many dislocation loops, which are not preserved after the load is removed, take part in the energy dissipation.

One more method of investigating the dynamic mobility of dislocations entails the measurement of the ranges of the lines and of the slip bands under the influence of a stress pulse of given duration [112-115]. However, the reduction of the results obtained by this method calls for special caution, since the leading dislocation is acted upon not only by an external stress but also by high internal stresses due to the other pile-up dislocations. As a result, the leading dislocation running away from the dislocations that pursue it, moves in an alternating field of stresses and consequently with variable velocity. This question was considered in detail in an experimental paper by Zaĭtsev and Nadgornyĭ [116].

Recently, after Indenbom and Éstrin [117] have demonstrated that not only the time of travel of the dislocations between barriers, but also the time of the thermal-fluctuation surmounting of the barriers is proportional to

TABLE I. Measured values of the drag coefficient B (millipoise)

Crystal / Method	LiF	NaCl	KCl	KBr	Cu	Al	Pb	Zn		Nb
								Basal	Pyramidal	
I	0,7 [71]	0,2 [72] 0,3 [82]	0,76 [3]	2,0 [74]	0,7 [76] 0,17 [81] 0,21 [83]	0,26 [79] 0,19 [84]		0,40 [80]	2,5 [75] 1,0 [86]	
II	1,3 [19] 0,34 [91]	0,16 [96] 0,11 [103]	0, 35 [19] 0,48 [104]	1,7 [95]	0,1— —0,8 [88] 0,65 [89] 0,12 [19] 0,2 [97] 0,85 [98]	2,0 [92] 1,7 [93] 3,1 [98]	0,37 [94]			0,17 [100]
III						0,25 [108] 0,19— —0,3 [110] 0,6 [114]	0,34 [113]			

Measurement methods: I—mobility of individual dislocations; II—amplitude-independent internal friction; III—experiments on high-speed deformation of the crystals and mobility of the slip bands.

the dynamic dragging coefficient, new possibilities were noted of analyzing the dynamic dragging by using macroscopic data, for example by using the influence of the superconducting transition on the plastic properties of metals.

Table I gives a summary of the main results of the measurement of the dynamic dragging of dislocations in various crystals, obtained by the methods listed above. All the data pertain to room temperature. The relatively large scatter of the values of B measured in different experiments on crystals of the same type may be due both to the large level of the experimental errors (especially in the indirect methods) and to the physical nonequivalence of the tested samples—for example, having different dislocation—"scaffold" densities—and this, according to [70] (see Sec. d of Chap. 3) should affect the level of the dynamic dragging of the dislocations.

The temperature dependence of the dragging coefficient B was first measured in copper by Alers and Thompson [88] by determining the amplitude-independent internal friction. Analogous measurements were subsequently carried out by various workers [19,90,91,93,94,97-105,112-114] on a number of other crystals, mainly by indirect methods and on the basis of the mobility of the slip bands. Unfortunately, these data were not very reliable since, as a rule, the temperature dependence of the dislocation density was not monitored during the reduction of the experimental data, and the principal difficulties of extracting information from experiments on the dynamics of slip bands were already mentioned by us before. In this connection, a particularly important role was played by recent measurements of the B(T) dependence in different crystals by direct methods [79-86]. In principle, the explanation of the functional B(T) dependence is a more serious criterion of the correctness and completeness of the theory than a simple comparison of a numerical estimate of B for one temperature. In particular, it turned out that allowance for the phonon wind, while giving the correct order of the dynamic drag, does not make it possible to explain the experimentally observed temperature dependence of B. It is precisely this circumstance which stimulated the search for new dissipation channels and revealed a specific role of the "slow" phonons.

As shown in Sec. e of Chap. 2, at not too low temperatures ($T \gtrsim \Theta/10$), any dislocation dragging is limited by

FIG. 9

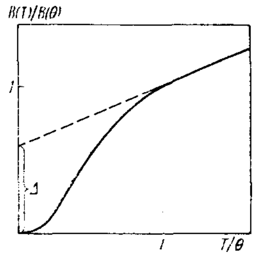


FIG. 10

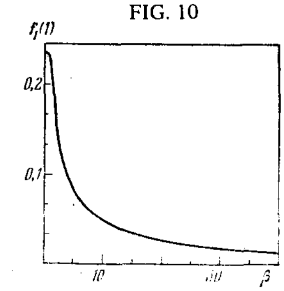


FIG. 9. Determination of the parameter Δ from the experimental B(T) curve.

FIG. 10. Plot of $f_1(1)$ against the parameter β.

a superposition of three effects: thermoelastic losses, phonon wind, and relaxation of "slow" phonons. The thermoelastic losses may turn out to be significant only at high temperatures $T \gg \Theta$, owing to their quadratic dependence on the temperature (2.32). Unfortunately, at the present time there are no reliable data on the high-temperature dynamic mobility of the dislocations. To be sure, Gektina and Lavrent'ev [86] in a recent measurement of the B(T) dependence for pyramidal dislocations in zinc observed, at high temperatures, deviations from linearity in the temperature dependence of B, which may indicate a contribution of the thermoelastic component. However, the excessively high level of the experimental errors casts doubts on the reality of these deviations.

At low temperatures ($T \lesssim \Theta/10$), where the flutter effects should predominate ($B \propto T^3$), and in metals, where electron scattering should prevail [118-121] (B = const) (see the note added in proof at the end of the article), the temperature dependence of B has also not been investigated thoroughly in experiment. There are only a number of partly contradictory data, obtained by an indirect method [88,90,97,98,101,105] and data on the mobility of the slip bands [112-114]. Notice should also be taken of a recent paper by Vreel and Jassby [85], who investigated the low-temperature mobility of individual dislocations in copper.

In the temperature region in which the bulk of the experimental material was obtained, the phonon slowing down of dislocations is determined by the last two terms of formula (2.30), which correspond to the contribution of the phonon wind and to the relaxation of the "slow" phonons:

$$B = \left[4 + \left(\frac{|n|}{G} - 6\right)^2\right] \frac{h}{b^3} \left(\frac{k_D b}{2\pi}\right)^5 \left[f_1\left(\frac{T}{\Theta}\right) + \lambda_\Theta \frac{\Theta}{T} f_2\left(\frac{T}{\Theta_S}\right)\right]. \quad (4.7)$$

In writing down (4.7), we took into consideration (2.21). The functions $f_1(x)$ and $f_2(x)$ are given by (2.24) and (2.31), while λ_Θ and Θ_S are phenomenological parameters to be determined from experiment. The theoretical determination of λ_Θ and Θ_S calls for knowledge of the real phonon spectra and of the real distortion field near the dislocation core. In any case, it is obvious that owing to the anisotropy of the phonon spectra each slip system for a given type of dislocation should generally speaking be characterized by its own parameters λ_Θ and Θ_S. We put throughout, for simplicity, $\Theta_S = \Theta$.

It is convenient to compare the temperature dependence of the drag coefficient with experiment in terms of the dimensionless coordinates $B(T)/B(\Theta)$ and T/Θ,

$$\frac{B(T)}{B(\Theta)} = \frac{f_1(T/\Theta)}{f_1(1)} (1 - \Delta \cdot f_2(1)) + \Delta \frac{\Theta}{T} f_2\left(\frac{T}{\Theta}\right), \quad \lambda_\Theta = \frac{\Delta f_1(1)}{1 - \Delta f_2(1)}, \quad (4.8)$$

where Δ is a dimensionless parameter determined from experiment by extrapolating to zero the temperature of the high-temperature asymptotic form of $B(T)/B(\Theta)$, which is usually linear in the temperature (Fig. 9); $f_2(1) \approx 0.92$, the numerical value of $f_1(1)$ being determined by the value of the parameter $\beta = 2k_D r_0$ (Fig. 10). To make the use of formulas (4.7) and (4.8) convenient, Fig. 11 shows plots of the functions $f_1(x)/f_1(1)$, $f_2(x)$, and $(1/x)f_2(x)$. It should be borne in mind here that in the scale used by us the curve $f_1(x)/f_1(1)$ is practically insensitive to the value of the parameter β. The reduction of the experimental data should begin with the extraction of the parameter β from the temperature dependence of B (Fig. 9). Then, choosing for r_0 a reasonable value (for example [4], $r_0 \approx 3b$), we calculate β, then $f_1(1)$, and finally the sought parameter λ_Θ. Knowing λ_Θ, we can use formula (4.7) to obtain the absolute value of B at any temperature. It must be borne in mind, however, that the factor Θ/T of the function $f_2(T/\Theta)$ in (4.7) and (4.8) is an approximate notation for the quantity l/l_Θ, which is valid only at temperatures that are not too low (see footnote 4). A better estimate is $l/l_\Theta \approx [C(\Theta)/C(T)]\kappa(T)/\kappa(\Theta)$, and in those cases when there are no experimental data on the temperature dependence of the lattice thermal conductivity κ (for example, for metals) we can assume, on the basis of the theoretical relation $\kappa \propto 1/T$, that $l/l_\Theta \approx [C(\Theta)/C(T)]\Theta/T$. Incidentally, the reduction of the experimental data by formulas (4.7) and (4.8) in their literal form, which is given below, seems to indicate that the approximation $l/l_\Theta \approx \Theta/T$ is perfectly satisfactory in many cases.

Figures 12—14 show the results of a comparison of the temperature dependence of $B(T)/B(\Theta)$ as given by formula (4.8) with the experimental points for a number of crystals. Table II gives the values of the parameters Θ and Δ, used to plot the theoretical curves. Unfortunately, the possibility of comparing the absolute values of B, calculated in accordance with formula (4.7), with the measured values, are limited to a small number of crystals, for which the values of the Murnaghan modulus n are known. Of the five crystals listed in Table II, the modulus n was measured only for copper. According to [122], $n/G \approx -33$ for copper. The corresponding estimates of the dragging coefficient for copper, in accordance with formula (4.7) yields as T = 300°K a value 1.6×10^{-12} poise, which agrees with the measured values (see Table I). It should be noted that this agreement is due to allowance for the relaxation of the elastic field of the dislocation near the nucleus, i.e., to the replacement of the function (2.20) by (2.24). An analogous estimate of B for copper, carried out in [25] in accordance with formulas of the type (2.19) and (2.20), has led to values that exceed the observed quantities by more than one order of magnitude.

FIG. 11

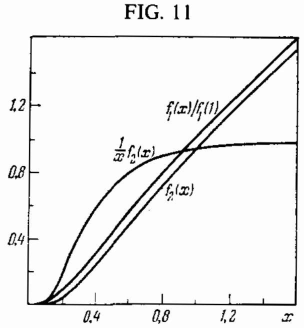


FIG. 12

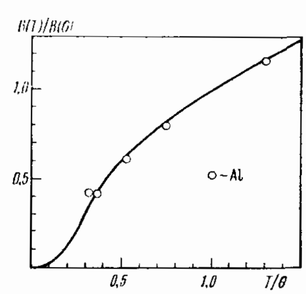


FIG. 11. Plots of the functions $f_1(x)/f_1(1)$, $f_2(x)$, and $f_2(x)/x$.
FIG. 12. Comparison of the temperature dependence of B(T) by formula (4.8) with the experimental points for an aluminum crystal (from the data of [79]).

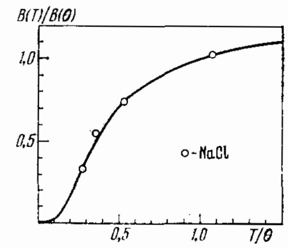


FIG. 13. Comparison of the temperature dependence of B(T) as given by formula (4.8) with the experimental points for an NaCl crystal (from the data of [82]).

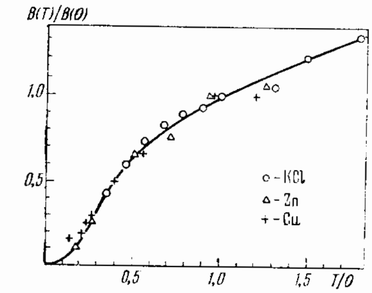


FIG. 14. Comparison of the temperature dependence of B(T) as given by formula (4.2) with the experimental points for the crystals Zn, Cu (from the data of [83], and KCl (from the data of [84]).

On the whole, it can be stated that there is good agreement between the theoretical and experimental results on the magnitude and temperature dependence of B(T). It is still necessary to investigate experimentally in detail the low-temperature part of this dependence, so as to reveal the contribution made to the dragging of the dislocations by the flutter effect (in non-metallic crystals) and by the electrons (in metals).

Another problem of interest, both theoretically and experimentally, is the influence of impurities on the dynamic dragging of dislocations. Unfortunately, only uncoordinated data are available so far on the influence of impurities. According to [78], the dynamic mobility of pyramidal dislocations in Zn is practically independent of the impurity concentration. On the other hand, it was shown in [3,82] that in the alkali-halide crystals KCl and NaCl one can discern the effect (albeit not a strong one) of impurities on the damping constant B.

Using the internal-friction method, Kaneda [101] separated the impurity contribution to the dragging coefficient B of copper doped with various concentrations of different impurities. It turned out that the impurity increment to B does not depend on the ratio of the masses of the impurity atoms to the matrix atom, is proportional to $C_0^{1/2}$ (C_0 is the impurity concentration), increases linearly with increasing impurity non-correspondence parameter ϵ_0, and reaches the same order of magnitude as the dragging coefficient in pure copper only at high impurity concentrations ($C_0 \sim 0.1$ at.% at $\epsilon_0 \sim 0.1$). The impurity

TABLE II

Crystal	Θ, °K	Δ	Crystal	Θ, °K	Δ
NaCl	280	0,91	Zn (basic)	240	0,61
KCl	270	0,61			
Cu	310	0,61	Al	230	0,49

influence can be traced also at low temperatures and cannot be attributed to the diffusion mechanisms, which predict an exponential decrease of the effect with decreasing temperature. The other mechanisms mentioned in Sec. d of Chap. 3 can likewise not explain the empirical relations obtained in [101].

It appears that to solve this problem it will be necessary to carry out new comprehensive experimental investigations of the dynamic mobility of dislocations in various crystals, as functions of the content and type of impurity at different temperatures, and also to undertake new theoretical exploratory studies. In addition, new investigations are needed to reveal the role of dislocation "scaffold" in dynamic losses. Work in this direction is already underway. One must mention the recent experimental paper by Gektina, Lavrent'ev, and Startsev [123], who have traced, in particular, the dependence of viscous slowing down of dislocations in zinc on the density of the "scaffold" dislocations.

Most phonon dissipation mechanisms lead, as we have seen, to viscous slowing down of the dislocations. An exception is radiation friction in the Peierls relief and a number of impurity effects characterized by the nonlinear dependence of the deceleration on the velocity. These exceptions, of interest in themselves, are however difficult to observe. Indeed, as shown above, radiation losses due to configuration oscillations of the dislocation field can be separated against the background of phonon-scattering processes only at near-sonic dislocation velocities and at low temperatures. Dynamic effects of the "dry friction" type would be observable only in pure crystals with a relatively high Peierls relief at dislocation velocities near the lower limit of the dynamic-velocity region, when the kinetic energy of the dislocation becomes comparable with the Peierls energy. To reveal these effects it is necessary to perform several experiments which apparently are still a matter for the future.

5. CONCLUSION

At the present time there is no longer any doubt concerning the role played by dynamic dragging in the mobility of fast dislocations, vibrational motion of dislocation segments between pinning centers, and processes of thermal fluctuation surmounting of local barriers by dislocation segments. For many crystals, experiment yields rather reliable data on the dragging coefficient and its dependence on the temperature, so that the theory can be verified. The main debatable problems in the theory of dynamic dragging of dislocations, with respect to the order of magnitude, region of applicability, and relative role of different mechanisms of phonon dragging, have been clarified.

A unified analysis of the various dissipative-loss channels points to a general hierarchy of the mechanisms of dynamic dragging of dislocations (Table III). The main dragging mechanism is usually phonon wind, which produces a linear temperature dependence of B at temperatures exceeding the Debye temperature and a relation $B \propto T^5$ at $T \ll \Theta$. The flutter is much less significant at $T \gtrsim \Theta$, but at low temperatures its contribution begins to prevail over the contribution of the phonon wind, owing to the more abrupt decrease (like T^3) with temperature. The phonon viscosity, in contrast to Mason's estimates, does not play any significant role in dislocation dragging, but an analogous effect, the relaxation of "slow" phonons, does make a noticeable contribution to the dynamic dragging at high temperatures. Since

TABLE III

Slowing-down mechanism	Temperature dependence of effect		B/B_w $(T \sim \Theta)$
	$T \ll \Theta$	$T \gtrsim \Theta$	
Phonon wind	T^5	T	1
Flutter effect	T^3	T	$\sim 10^{-1}$
Relaxation of "slow" phonons	$\frac{\tau}{T} e^{-\Theta/T}$	const	~ 1
Contribution of optical modes	$\frac{\tau}{T} e^{-\Theta_0/T}$	$AT+B$	$\sim \left(\frac{\Theta}{\Theta_0}\right)^4 < 1$
Thermoelastic losses	T^4/τ	T^2	$\sim 10^{-1}$
Phonon viscosity	T^4/τ	T^2	$\sim \left(\frac{v}{c}\frac{b}{l}\right)^2 \ll 1$
Raman scattering	T^5	T	$\sim 10^3 \frac{\sigma_P}{G}$
Radiation friction			$\frac{\sigma_P b}{B_w v} \operatorname{cth} \frac{\sigma_P b}{B_w v} - 1$

this contribution tends little on the temperature at $T \gtrsim \Theta$, it exerts a decisive influence on the character of the temperature dependence of the slowing down. The thermoelastic losses are usually insignificant, but in individual cases they can become manifest at high temperatures, owing to the quadratic dependence of the effect on the temperature. The contribution of the optical phonons, as a rule, is also small and can appear only under special conditions (soft modes that are produced in phase transitions, etc.).

The Peierls relief can appear both in the radiative dragging of the dislocations and in the Raman scattering of the phonons in the oscillating field of a dislocation. The former effect may in principle turn out to be responsible for the phenomenon of dry friction in the case of slowly moving dislocations (and kinks). The latter effect is significant in comparison with the phonon wind only for crystals with a high Peierls relief ($\sigma_P/G \gtrsim 10^3$).

Further refinement of the theoretical estimates calls for a complete allowance for the anisotropy of the crystal and for a consideration of concrete phonon spectra. Such calculations are of particular interest in the case when one can expect a noticeable anisotropy of the effect and a sharp temperature dependence of the effect (for example, near phase transitions). The launching of the corresponding experimental research has by now become quite timely.

The problem of the influence of irradiation and of doping on the dynamic dragging of dislocations calls for additional theoretical and experimental research. From the presently available experimental data it is impossible to determine the causes of the observed disparity between experiment and theory.

A number of predictions of the theory are of great physical interest, and undoubtedly are worthy of experimental verification. These include the effects of the dry friction, critical velocity, radiative dragging at near-sonic velocities, excitation of local and quasilocal oscillation of impurity atoms, dragging by soft modes, and low-temperature manifestation of the flutter contribution. In view of the latest progress in the experimental and experimental research of dynamic slowing down, there are grounds for hoping that in the nearest future considerable progress will be made with respect to all the aforementioned topics.

In this review we confined ourselves to an analysis of phonon mechanisms of dislocation dragging. The electronic mechanisms of slowing of dislocations in metals are almost as numerous and varied (see the recently published reveiw [124]) and the mechanisms of dynamic dragging of kinks (see the review [125]).

In conclusion, the author thanks I. M. Lifshitz and M. I. Kaganov for useful discussions and valuable remarks, and also V. I. Startsev, E. M. Nadgornyĭ, and F. F. Lavrent'ev for the opportunity of becoming acquainted with new experimental data.

Note added in proof. As shown in a recent paper by Al'shitz [126], in individual metals an important role may be played by the relaxation of electrons belonging to flattened Fermi-surface sections normal to the dislocation line. The corresponding dragging is proportional to the electric conductivity of the metal and increases with decreasing temperature.

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1) The attempt to solve this problem in [20] turned out to be incorrect, as pointed out in [25].

2) It is seen from (2.19) that at low temperatures we have $B \propto c^{-5}$. Usually the velocity of the transverse phonons (c_t) is much lower than the velocity of the longitudinal phonons (c_l) and therefore the contribution of the longitudinal phonons to the dragging is smaller by a factor of $2(c_l/c_t)^5$ than that of the transverse phonons.

3) It suffices in this case to know the asymptotic expressions presented above for $\eta_{ijkl}(q, \Omega_q)$ at large and small values of q, since the first integral depends on the upper limits only logarithmically, and the second is practically independent of the lower limit.

4) The phonon mean free path l can be estimated from the thermal diffusivity $\chi = \kappa/C$: $l \approx 3\chi/c = 3\kappa/cC$. Usually, starting with sufficiently low temperatures ($T \gtrsim \Theta/10$), the lattice thermal conductivity κ varies in proportion to T^{-1}, and the specific heat C depends little on the temperature. This makes it possible to describe the temperature dependence of l by the relation $l(T) \approx l_\Theta\, \Theta/T$ in the entire temperature region in which the "slow" phonons are not "frozen out."

5) The calculation in [31] was carried out in a somewhat rougher approximation, with the same $\Psi(y) = \text{const} \sim 1$ for screw and edge dislocations.

6) Recognizing that the foregoing estimate of the flutter effect was obtained without taking into account the specific contribution of the dislocation core, we compare expressions (3.2) and (3.3) with the formulas (2.19) and (2.20). The quantity T_0 which pertains to the region of low temperatures, where the problem is not very sensitive to the dimension of the dislocation core, is estimated more reliably than the quantity d_0.

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