Theory of Metal Weakening During a Superconducting Transition

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A theory is developed of metal weakening effects which had been observed experimentally in transitions to the superconducting state and which consist in a noticeable increase of plasticity of the metals. The case of not very large stresses is considered. For such stresses the plastic deformation rate is determined by thermally activated motion of the dislocations through a set of potential barriers produced in the crystal by various structural defects. It is shown that the mean time required for overcoming the barrier depends significantly on the electron retardation force of the dislocation. As a result a sharp change of the force during the superconducting transition leads to a change in the kinetics of the plastic deformation. Weakening effects are analyzed for various deformation conditions created experimentally, for active deformation at a constant rate, creep, or stress relaxation. The temperature dependence of the effects is investigated and compared with the experimental data.

1. INTRODUCTION

MANY experimental data published in recent years offer evidence of a peculiar weakening, consisting of a noticeable increase in plasticity, occurring in metals going over into the superconducting state. This phenomenon was observed under a great variety of deformation conditions. In particular, it turned out that when a metal becomes superconducting its elastic limit can decrease by several dozen per cent[1], and an anomalous behavior can be observed with decreasing temperature[2]; the deformation stress required to maintain a constant rate of plastic flow also decreases by several per cent[3]. A particularly large effect was observed in creep experiments, viz., the transition of a metal into the superconducting state during the stage of steady-state creep is accompanied by a sharp increase in the creep rate by tens or even hundreds of times[4,5]. There are data indicating that the superconducting transition exerts an influence on the stress relaxation[6-8], on the mobility of the slip bands[9], and on many other plastic characteristics of a metal. We present here reference only to the first investigations in which the indicated effects were observed; at the present time each effect is the subject of several studies.

When an attempt is made to systematize the experimental results, the following characteristic features of the observed phenomena call attention to themselves. First, within the limits of the measurement accuracy, no influence of the superconducting transition on the process of elastic deformation was observed; noticeable effects appear only during the stage of well developed plastic deformations, when a displacement of a large number of dislocations takes place in the crystal. Second, in all cases without exception the superconducting transition is accompanied by an increase in the plasticity of the metal. Third and finally, a definite correlation is observed between the temperature dependence of the weakening effects and the temperature dependence of the energy gap of the superconductor. These distinguishing features of the weakening allow us to propose that they are based on an appreciable increase of dislocation mobility, which occurs when the metal becomes superconducting, and that this increase is closely connected with the realignment of the energy spectrum of the metal in such a transition.

The mechanisms whereby dislocations moving in a real crystal are slowed down are customarily subdivided at present into two groups, depending on their physical nature. Let us discuss the possible influence of the superconducting transition on each of these groups.

The first group includes slowing-down mechanisms due to the presence of different defects in the crystal, impurities and their clusters, other dislocations, etc., which are sources of potential barriers hindering the slip. Such defects play a double role. On one hand, they give rise to long-range elastic fields that produce in the crystal "ripples" of internal stresses with an amplitude σ_i that depends on the defect density; as a result, the slipping of the dislocations becomes possible only if the external stress σ applied to the dislocation exceeds σ_i. On the other hand, defects that fall in the slip plane produce barriers that are localized at atomic distances. These barriers, because of their small spatial dimension, are overcome by the dislocation even at a low level of internal stress, by thermal activation. Among the thermally surmountable barriers are also the Peierls-Nabarro barriers, which are due to the discrete structure of the crystal. Thus the mobility of a dislocation at low stresses is determined by the rate of thermal-activation surmounting of the indicated barriers.

The main influence is exerted on the process of thermal activation by the height of the barriers, so that the possibility of changing this height should be the first item under discussion. Such a possibility exists because in general the elastic properties of a crystal lattice are sensitive to the superconducting transition. It has been established, however, both theoretically and experimentally[10], that the change in the properties of the crystal lattice is exceedingly small (the relative change of the elastic moduli is $\sim 10^{-5}$). Very simple estimate show that the influence of such small changes cannot lead to the appearance of noticeable effects in such rather "crude" measurements as the usual mechanical tests of metals.

The second group of deceleration mechanisms includes different dissipative processes which occur in the elastic field of the moving dislocation. These include phonon and electron viscosities, thermoelastic losses, etc. The part of the deceleration force generated by them becomes predominant only at very large dislocation velocities ($\sim 10^4 - 10^5$ cm/sec), and it is therefore customarily assumed that the role of these mechanisms is negligible at low velocities. The dissipative properties of the metals at low temperatures, as is well known, are determined by the absorbing ability of the conduction electrons. The dislocation deceleration force component due to the electron viscosity was calculated for a normal metal by Kravchenko[11] and Holstein[12] and turned out to be proportional to the dislocation velocity V:

$$f_n = B_n V, \tag{1}$$

where B_n is the damping coefficient, the magnitude of which does not depend on the temperature and is of the order of 10^{-5} g/cm-sec for typical metals. This part of the deceleration force should decrease strongly during the superconducting transition, since the absorbing ability of the conduction electrons is strongly decreased when a gap Δ appears in their energy spectrum.

The electron deceleration force calculated in[13,14] in a superconductor turned out to be a complicated nonlinear function of the temperature and of the dislocation velocity. A linear dependence on the velocity occurs only at relatively low velocities $V \ll \Delta/p_F$, T/p_F (p_F is the Fermi momentum), and unlike in a normal metal the damping coefficient depends here strongly on the temperature:

$$f_s = B_s V, \quad B_s = 2B_n / (1 + e^{\Delta/T}). \tag{2}$$

An analysis of the temperature dependence of the jump of the electron-deceleration force of the dislocation on going from the normal to the superconducting state shows[13] that this dependence is close to the dependence $\Delta(T)$, although they are not exactly equal.

However, when an attempt is made to apply the results of[13,14] to an explanation of the weakening phenomena, it should be borne in mind that a direct connection between the jump in the force of electron deceleration and the change of the characteristics of the plastic deformation exists only in those few cases when the dislocation moves during the course of deformation with high velocities, at which their kinetic energy exceeds the heights of the potential barriers. Such velocities are realized, for example, in the frontal part of a slip band propagating under the action of large pulsed loads[9], and in the case of active deformation apparently at large deformation rates and high stresses. As to the majority of the observed weakening effects, they are obtained under deformation conditions (creep, stress relaxation, initial stage of active deformation) when the plastic deformation is more readily effected by dislocations thermally overcoming the potential barriers, and therefore the magnitude of the weakening should be determined by the influence of the superconducting transition on this process.

Since, as already noted, there are no grounds for assuming that the heights of the barriers can change noticeably in the superconducting transition, it remains to propose that the thermal-activation process is sensitive to changes in the electronic damping of the dislocations. Such a proposal is perfectly natural and agrees with the general principles of fluctuation theory, according to which the thermal fluctuations are closely connected with the dissipative properties of the medium.

It is shown in the present paper that electronic damping in metals at low temperatures exerts a very strong influence on the thermal-activation motion of the dislocation through barriers, and that its variation during the superconducting transition explains almost all the presently existing experimental results on weakening. In the second section of the paper we investigate the thermal motion of a dislocation segment pinned by defects and calculated with the average frequency of detachment from an individual defect. For sufficiently long dislocations, this frequency turns out to be inversely proportional to the damping coefficient, whereas the thermal motion of short segments of dislocations is in practice not sensitive to the decleration force. In the third section, with the simplest dislocation model of plastic deformation as an example, the influence of an abrupt change in the electronic damping at the instant of the superconducting transition on the kinetic deformation of the metal is analyzed. Different damping conditions realized in experiments are considered separately, namely, active deformation, creep, and stress relaxation. The weakening of superconductors that go over from the normal into the intermediate or mixed state is also analyzed.

2. THERMAL MOTION OF DISLOCATIONS

The elementary act of plastic deformation in thermal-activation motion of dislocations through an aggregate of potential barriers is shown schematically in Fig. 1. If a dislocation moving under the influence of an external stress σ has acquired a configuration AC_1B, then following a certain time interval it assumes, after experiencing a thermal fluctuation of the proper amount, the position AC_2B, after which the stress σ moves it into the position AC_3B. An exact calculation of the average time required for the dislocation to overcome the barrier is a complicated problem, which

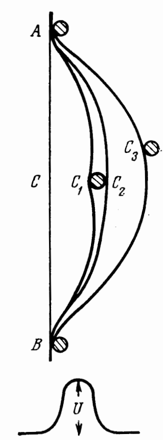

FIG. 1

has been considered by many authors, but has not been solved rigorously to date. In most cases, the methods of the theory of absolute reaction rates were used in the calculation of this time, but no account was taken at all of the dissipative properties of the crystal[1)]. In the present paper we also confine ourselves to the approximation of the theory of reaction rates, and attempt to clarify the role played in this approximation by the dissipative properties of the crystal among the main factors which determine the process of thermal activation of the dislocations. Of course, the assumption made above concerning the influence of the dissipative properties on the crystal on the process of thermal-activation displacement of the dislocation lie outside the scope of the reaction-rate theory and should be taken into account in any more rigorous approach.

The lower part of Fig. 1 shows schematically the energy of interaction of the dislocation with the defect; region I in this figure corresponds to the position AC_1B, region III to the position AC_2B, and the intermediate region II corresponds to the "excited" state, in which the dislocation is located at the vertex of the potential barrier. The rate of the stationary process of the thermal-activation motion of the dislocation through the barriers is determined in the approximation of the reaction-rate theory by the relation

$$w = \nu e^{-H(\sigma)/T}, \tag{3}$$

where w is the average frequency for overcoming the barrier, $H(\sigma)$ is the activation energy with allowance for the external stress σ, T is the crystal temperature in energy units, and ν is the "collision" frequency of the dislocation with the barrier in its thermal motion in region I. When formula (3) is applied, it is customary to pay principal attention to the analysis of $H(\sigma)$, and the factor ν is set in most cases equal to the frequency of the natural oscillations of the dislocation segment[16]. It will be shown below that the value of ν in (3) depends, generally speaking, on the deceleration force acting on the dislocation, and in a number of cases such a dependence is quite significant.

Proceeding to the calculation of the frequency ν, we note first that during the time of stay in region I the dislocation is in thermal equilibrium with the crystal, and therefore ν is none other than the characteristic frequency of the thermal oscillations of the dislocation, which can be obtained on the basis of the theory of equilibrium thermal fluctuations. It is easy to show that the order of magnitude of this frequency does not depend on the presence of a barrier that limits the dislocation motion on one side, or on the presence of the homogeneous external stress σ; the role of the latter reduces to a certain renormalization of the linear tension of the dislocation[17]. Therefore the frequency ν is best estimated by considering the free thermal oscillations of the dislocation segment AB about the straight line joining its ends.

We denote the displacement of the dislocation-segment element from the straight line AB by u(x, t), choosing the x axis along this straight line with origin at the point A; we assume that the function u(x, t) satisfies the condition that the points A and B are fixed: $u(0, t) = u(L, t) = 0$, where L is the length of the segment. At random thermal fluctuations, the average displacement $\overline{u(x, t)}$ is zero, and therefore the thermal motion of the dislocation should be characterized by the mean-squared displacement $\overline{u^2(x, t)}$. The thermodynamic equilibrium value of this quantity, $\overline{u_0^2(x)}$, was calculated earlier by Leibfried[18]. It is, naturally, independent of the time and in the classical limit its value for the center of the segment $(x = L/2)$ is

$$\overline{u_0^2(\tfrac{1}{2}L)} \equiv \overline{u_{0c}^2} = 2T / LM\omega_0^2, \quad \omega_0^2 = \pi^2 C / L^2 M, \tag{4}$$

where M and C are respectively the linear density of the effective mass and the coefficients of linear tension, and ω_0 is the natural frequency of the segment.

If we calculate the mean-squared segment displacement $\overline{u^2(x, t)}$ from the initial position $u(x, 0) = 0$, then this quantity should tend in the course of time to an equilibrium value $\overline{u_0^2(x)}$; for the central point $x = L/2$ we have

$$\lim_{t\to\infty} \overline{u_c^2(t)} = \frac{2T}{LM\omega_0^2}, \quad \overline{u_c^2(t)} \equiv \overline{u^2(\tfrac{1}{2}L, t)}. \tag{5}$$

The mean value of the period of the thermal oscillations which is of interest to us, can be defined as the characteristic time during which the mean-squared segment displacement $\overline{u^2(x, t)}$ from the initial position $u(x, 0) = 0$ reaches its equilibrium value $\overline{u_0^2(x)}$. To find this time, it suffices to trace the time dependence of the displacement $\overline{u_c^2(t)}$ of the segment center.

The value of $\overline{u_c^2(t)}$ can be calculated on the basis of the well known "string" model of the dislocation, according to which the oscillations of the segment are determined by the equation

$$M\ddot{u} - Cu'' + BV = 0, \tag{6}$$

where $\dot{u} \equiv \partial u/\partial t$, $u' \equiv \partial u/\partial x$, and $V = \dot{u}$. The last term in (6) describes the force of the electronic slowing down of the dislocation, and therefore the damping coefficient B can take on the value B_n or B_s, depending on the state of the metal. The choice of a linear velocity dependence of the decelerating force in the case of a superconductor is justified by the small thermal velocity of the dislocation. The mean-square fluctuation of the velocity can be obtained on the basis of (4), namely, $\overline{V^2} \sim u_0^2\omega_0^2 = 2T/LM$; comparing the obtained estimate with the condition for the applicability of formula (2), we see that in view of this condition $(\sqrt{\overline{V^2}} \sim \sqrt{2T/LM} \ll \Delta(T)/p_F, T/p_F)$ is satisfied in almost the entire temperature interval, with the exception of small vicinities of absolute zero $T = 0°K$ and the superconducting transition temperature $T = T_c$.

Expanding u(x, t) in a Fourier series, we can represent the motion of the dislocation segment in the form of oscillations of a set of harmonic oscillators,

$$u(x, t) = \sum_n a_n(t) \sin\frac{\pi n}{L}x, \quad \overline{u_c^2(t)} = \sum_n \overline{a_n^2(t)} \sin^2\frac{\pi n}{2}.$$

As shown by Leibfried[18], the mean-squared fluctuation is determined mainly by the first term of this sum

[1)]An exception is a recent paper by Indenbom and Estrin[15], who noted the connection between the mobility of the dislocation as it moves in front of the barrier and the average time necessary to overcome the barrier. Their paper discusses also the limits of applicability of the approximations of the reaction-rate theory to this problem.

($n = 1$). Consequently, $u_c(t) \approx a_1(t)$ and satisfies the equation

$$\ddot{u}_c + \omega_0^2 u_c + \gamma \dot{u}_c = 0, \tag{7}$$

where $\gamma = V/M$ is the damping coefficient of the free oscillations of the dislocation.

According to the general fluctuation theory[19], spontaneous fluctuations of the quantity $u_c(t)$ can be considered as the result of the action of a certain fictitious random force $f(t)$. Substituting this force into the right-hand side of (7) and differentiating it with respect to time, we obtain an equation for the velocity $V_c = \dot{u}_c$:

$$\ddot{V}_c + \omega_0^2 V_c + \gamma \dot{V}_c = \dot{f}(t). \tag{8}$$

We represent the solution of Eq. (8) in the form of a Fourier integral

$$V_c(t) = \int_{-\infty}^{\infty} d\omega [-i\omega\alpha(\omega) f^\omega] e^{-i\omega t},$$

where f^ω is the spectral component of the random force and $\alpha(\omega)$ is the generalized susceptibility of the dislocation

$$\alpha(\omega) = [\omega_0^2 - \omega^2 - i\gamma\omega]^{-1}. \tag{9}$$

After this, we obtain next an expression for the quantity $\overline{u_c^2(t)}$ of interest to us:

$$\overline{u_c^2(t)} = \left[\int_0^t V_c(t)\,dt\right]^2 = \int_{-\infty}^{\infty} d\omega \int_{-\infty}^{\infty} d\omega' \alpha(\omega)\alpha(\omega')(1 - e^{-i\omega t})(1 - e^{-i\omega' t})\overline{f^\omega f^{\omega'}}. \tag{10}$$

The mean-squared fluctuation of the random force is determined by the expression[19]

$$\overline{f^\omega f^{\omega'}} = A \frac{\hbar \operatorname{Im} \alpha(\omega)}{|\alpha(\omega)|^2} \operatorname{cth} \frac{\hbar\omega}{2T} \delta(\omega + \omega'), \tag{11}$$

where A is a normalization constant that can be determined from the condition (5).

Substituting (11) in (10) and taking into account the obvious property of the generalized susceptibility $\alpha(-\omega) = \alpha^*(\omega)$, we obtain

$$\overline{u_c^2(t)} = \frac{A\hbar}{2\pi} \int_{-\infty}^{\infty} d\omega \operatorname{Im} \alpha(\omega)(2 - e^{i\omega t} - e^{-i\omega t}) \operatorname{cth} \frac{\hbar\omega}{2T}.$$

This integral is relatively easy to calculate with the aid of residue theory, and the final expression for $\overline{u_c^2(t)}$ is

$$\overline{u_c^2(t)} = \frac{A\hbar}{\sqrt{\gamma^2 - 4\omega_0^2}} \left[(1 - e^{-\lambda^{(-)}t}) \operatorname{ctg} \frac{\hbar\lambda^{(-)}}{2T} - (1 - e^{-\lambda^{(+)}t}) \operatorname{ctg} \frac{\hbar\lambda^{(+)}}{2T}\right],$$
$$\lambda^{(\pm)} = \tfrac{1}{2}(\gamma \pm \sqrt{\gamma^2 - 4\omega_0^2}). \tag{12}$$

Of practical interest to us is the classical limit of expression (14), which occurs under the condition $\hbar\lambda^{(\pm)} \ll T$:

$$\overline{u_c^2(t)} = \frac{2AT}{\sqrt{\gamma^2 - 4\omega_0^2}} \left(\frac{1 - e^{-\lambda^{(-)}t}}{\lambda^{(-)}} - \frac{1 - e^{-\lambda^{(+)}t}}{\lambda^{(+)}}\right). \tag{13}$$

Substituting (13) in (5), we obtain the value of the normalization constant $A = 1/LM$.

Formula (13) enables us to establish one important feature of the thermal motion of the dislocation segment. Although in the classical limit the equilibrium value of the displacement fluctuation does not depend on the dissipative properties of the crystal, as should be the case in accordance with the theory of classical fluctuations[19], the time necessary to reach such an equilibrium value from a specified initial position (a characteristic period of the thermal motion of the segment) turns out to be quite sensitive to these properties. In order to make this statement more obvious, let us investigate the time dependence of the quantity $\overline{u_c^2(t)}$ for three limiting ratios of the damping coefficient γ to the natural frequency of the segment ω_0:

$$\overline{u_c^2(t)} = \begin{cases} \dfrac{2T}{LM\omega_0^2}(1 - \cos\omega_0 t), & \gamma \ll \omega_0; \\ \dfrac{2T}{LM\omega_0^2}[1 - (1 + \omega_0 t)e^{-\omega_0 t}], & \gamma = 2\omega_0; \\ \dfrac{2T}{LM\omega_0^2}\left(1 - \exp\left\{-\dfrac{\omega_0^2}{\gamma}t\right\}\right), & \gamma \gg \omega_0,\ t > \dfrac{1}{\gamma}. \end{cases} \tag{14}$$

From this we see clearly the influence exerted by the dissipative properties of the crystal, which are characterized by the coefficient γ, on the thermal motion of the dislocation. At relatively small damping $\gamma \lesssim 2\omega_0$, the characteristic period of the thermal oscillations coincides in order of magnitude with the period $1/\omega_0$ of the natural oscillations of the segment and is not sensitive in practice to the value of γ. But at a large relative damping $\gamma > 2\omega_0$ this period is of the order of $\gamma/\omega_0^2 \gg 1/\omega_0$, i.e., it increases sharply and becomes proportional to the coefficient γ.

The criterion that determines the cases of weak and strong damping should be discussed in greater detail. In a normal metal, the damping coefficient $\gamma = B_n/M$ is a constant quantity, and therefore, if we introduce the critical length L_0 of the segment, defined by the relation

$$L_0 = 2\pi\sqrt{MC}/B_n, \tag{15}$$

then the damping will be weak for segments of length $L < L_0$ and strong for the longer segments with $L > L_0$. For typical metals with $B_n \sim 10^{-5}$ g/cm-sec, the length L_0 is of the order of 10^{-5} cm.

In a superconductor, according to (2), the damping coefficient $\gamma = B_s/M$ decreases rapidly with decreasing temperature. It is therefore convenient to use as the criterion separating the cases of strong and weak damping the temperature T_0, defined by the equation $B_s(T) = 2M\omega_0$; then we have weak damping at $T < T_0$ and strong damping at $T > T_0$. The temperature T_0 depends, naturally, on the segment length L; the equation that determines this length can be written with the aid of formula (2) in the form

$$\frac{\Delta(T)}{T} = \ln\left(\frac{2L}{L_0} - 1\right). \tag{16}$$

Numerical solution of (16) yields $T_0(10^3 L_0) \approx T_c/5$, $T_0(10^2 L_0) \approx T_c/4$, and $T_0(10L_0) \approx T_c/2$ (T_c is the critical temperature of the superconducting transition).

The foregoing analysis of the thermal motion of a dissociation segment makes yields the following estimate for the dependence of the pre-exponential factor

ν in (3) on the electronic slowing-down coefficient

$$\nu \approx \begin{cases} \omega_0, & B \lesssim B_n L_0/L; \\ M\omega_0^2/B, & B > B_n L_0/L. \end{cases} \tag{17}$$

We note that a similar result can be obtained by considering a somewhat different scheme of detaching the dislocation from the defect. If the interaction of the dislocation with the defect is sufficiently strong, then the dislocation element adjacent to the defect will be "firmly" secured (at the point C_1 in Fig. 1): in this case the detachment will occur when the angle between the straight line C_1B and the tangent to the dislocation line at the point C_1 reaches a certain critical value. The average frequency of such an event in the case of random thermal motion of the segment C_1B can also be determined with the aid of (3), but the frequency ν in this formula will now have the meaning of the frequency of variation of the indicated angle. Within the framework of the string model for dislocations, this frequency coincides with the frequency of motion of the center of the segment C_1B_1 and is consequently determined by formula (17), in which L must be taken to mean the length of the segment C_1B.

We have thus established that at a sufficiently large distance between the barriers the pre-exponential factor in (3) becomes proportional to the dislocation damping coefficient. The result confirms the assumption made in the introduction concerning the influence of the dissipative properties of the crystal on the thermally-activated motion of the dislocation. Any sharp change in the dislocation slowing-down mechanism in a sufficiently perfect crystal leads to an equally sharp change in the average time of the thermal activation, which in final analysis will exert a substantial influence on the kinetics of the plastic deformation.

3. INFLUENCE OF SUPERCONDUCTING TRANSITION ON THE PLASTIC DEFORMATION OF A METAL

The transition from the normal to the superconducting state (ns) and the inverse transition (sn) are realized experimentally either by turning on and off a superconductivity-destroying magnetic field, or by gradually decreasing the temperature. In the former case the coefficient of the electronic damping B changes practically instantaneously by an amount $B_n - B_s$, whereas in the latter case it changes in synchronism with the change of the temperature in accordance with formula (2). We shall analyze the influence of such changes on the deformation process using as an example the simplest dislocation model of plasticity.

Assume that the crystal contains one slip system with density N of mobile dislocations and an aggregate of thermally-surmountable barriers of one type with energy U. The slip of the dislocation after a single act of thermal activation (Fig. 1) is accompanied by plastic deformation bS, where b is the magnitude of the Burgers vector and S is the area of the figure AC_1BC_3. If the average distance between barriers is L/2, then the number of segments of the type AC_1B per unit volume is N/L, and the rate of plastic deformation $\dot{\epsilon}$ in the stationary regime is given by the expression

$$\dot{\varepsilon} = bSNw/L. \tag{18}$$

In writing down (18), we assume that the external stress σ is not very large, so that the time necessary to overcome barrier 1/w greatly exceeds the time of dislocation travel from the position AC_2B to the position AC_3B.

The activation energy $H(\sigma)$ takes in the simplest case the form[16]

$$H(\sigma) = U - v(\sigma - \sigma_i), \tag{19}$$

where v is the activation volume and σ_i is the internal-stress component acting in the slip plane. Substituting (3) in (18) and taking (19) into account, we have

$$\dot{\varepsilon} = \varepsilon_0 \nu \exp\left\{-\frac{U - v(\sigma - \sigma_i)}{T}\right\}, \quad \varepsilon_0 = \frac{bSN}{L}. \tag{20}$$

The strengthening process occurring during the time of the plastic deformation can be taken into account phenomenologically, by assuming that stresses σ_i increase in proportion to the strain[16]:

$$\sigma_i = \sigma_i^0 + k(\varepsilon - \varepsilon_u). \tag{21}$$

Here ϵ_u is the elastic strain, k is the strengthening coefficient, and σ_i^0 are the internal stresses at the instant of the start of plastic deformation.

Starting from formulas (20) and (21), let us consider separately different deformation conditions. We note immediately that in the case of relatively imperfect crystals ($L < L_0$), the frequency ν is practically insensitive to electronic slowing down of the dislocation, and the supercondition transition should not lead to a weakening of such crystals. We shall therefore consider from now on only the case $L > L_0$.

Active Deformation with Constant Velocity

In this case $\dot{\epsilon}$ = const and the deforming stress increases with time; integrating (20), we have

$$\sigma(t) = \sigma_i + \frac{U}{v} - \frac{T}{v}\ln\frac{\varepsilon_0\nu}{\dot{\varepsilon}} + k(\dot{\varepsilon}t - \varepsilon_u). \tag{22}$$

Let us consider first the time-independent part of the deforming stress—the yield point σ_k:

$$\sigma_k = \sigma_i + \frac{U}{v} - \frac{T}{v}\ln\frac{\varepsilon_0\nu}{\dot{\varepsilon}}. \tag{23}$$

On passing through the critical temperature T_c, the temperature dependence of this quantity reveals a number of singularities determined by the temperature behavior of the frequency ν.

As already noted above, in the case $L < L_0$ the $\sigma_k(T)$ dependence should retain its usual form on going through T_c:

$$\sigma_k(T) = \sigma_i + \frac{U}{v} - \frac{T}{v}\ln\frac{\varepsilon_0\omega_0}{\dot{\varepsilon}}, \quad L < L_0. \tag{24}$$

In more perfect crystals ($L > L_0$), the $\sigma_k(T)$ dependence becomes nonmonotonic near the temperature T_c:

$$\sigma_k(T)=\begin{cases}\sigma_i+\dfrac{U}{v}-\dfrac{T}{v}\ln\dfrac{M\varepsilon_0\omega_0^2}{\dot{\varepsilon}B_n}, & T>T_c\\[2mm] \sigma_i+\dfrac{U}{v}-\dfrac{T}{v}\ln\left[\dfrac{M\varepsilon_0\omega_0^2}{\dot{\varepsilon}B_n}\,\dfrac{1+e^{\Delta/T}}{2}\right], & T_c>T>T_0\\[2mm] \sigma_i+\dfrac{U}{v}-\dfrac{T}{v}\ln\dfrac{\varepsilon_0\omega_0}{\dot{\varepsilon}}, & T<T_0.\end{cases}\tag{25}$$

The temperature dependence of $\sigma_k(T)$ at different distances between the barriers is shown schematically in Fig. 2. In the experiments[2], they investigated the temperature dependence of a quantity close in its physical meaning to σ_k, namely the yield point σ_u (the difference between them is shown in Fig. 3). When the metal goes over into the superconducting state, σ_u decreases sharply, and then again begins to increase at lower temperatures, in qualitative agreement with the results obtained here.

If the superconducting transition occurs instantaneously (with the aid of a magnetic field), then a jumplike change of the frequency ν will occur at the instant of the transition and will lead to a jump-like change of the deforming stress. The jump of the deforming stress is equal to

$$\delta\sigma_{sn}=-\delta\sigma_{ns}=\frac{T}{v}\ln\frac{\nu_s}{\nu_n}$$

(ν_s and ν_n are the values of the frequency ν in the superconducting and normal states), and, depending on the temperature, its value is

$$\delta\sigma_{sn}=\begin{cases}\dfrac{T}{v}\ln\dfrac{1+e^{\Delta/T}}{2}, & T_c>T>T_0\\[2mm] \dfrac{T}{v}\ln\dfrac{B_n}{M\omega_0}, & T<T_0.\end{cases}\tag{26}$$

Consequently, the deformation curve for the case discussed here should have the form shown schematically in Fig. 3. The dashed line in the figure shows the deformation curve observed experimentally; certain deviations of the theoretical curve from the experimental one in the region of the yield point at the instants of the sn and ns transitions are apparently connected with the nonstationary character of the process of thermal activation motion of the dislocation at these instants. The temperature dependence of the quantity $\delta\sigma_{sn}$ is discussed later.

Creep

In the case of creep, σ = const and the strain ϵ increases with time. Integration of (20) with allowance for (21) yields

$$\varepsilon(t)-\varepsilon_u=\frac{T}{kv}\ln(\alpha\nu t+1),\quad \dot{\varepsilon}(t)=\alpha\nu T/kv(\alpha\nu t+1),$$
$$\alpha=\frac{kv\varepsilon_0}{T}\exp\left\{-\frac{U-v(\sigma-\sigma_i^0)}{T}\right\}.\tag{27}$$

This is the usual logarithmic creep, shown schematically by the lower curve of Fig. 4. The instantaneous transition of the crystal from the normal to the superconducting state ($t=t_0$) changes abruptly the course of the creep process. In general form, the course of the process can be described by the formula

$$\varepsilon(t)-\varepsilon_u=\begin{cases}\dfrac{T}{kv}\ln(\alpha\nu_n t+1), & t<t_0\\[2mm] \dfrac{T}{kv}\ln[\alpha\nu_s(t-t_0)+\alpha\nu_n t_0+1], & t>t_0.\end{cases}\tag{28}$$

According to this formula, the creep rate $\dot{\epsilon}(t)$ at the instant of transition increases jumpwise by a certain amount $\delta\dot{\epsilon}_{ns}$, as a result of which an additional strain increment $\delta\epsilon_{ns}$ is obtained (Fig. 4). The jump in the rate of deformation and the elongation of the sample after the ns transition take the form

$$\delta\dot{\varepsilon}_{ns}=\dot{\varepsilon}_n(t_0)\left(\frac{\nu_s}{\nu_n}-1\right),\quad \delta\varepsilon_{ns}=\frac{T}{kv}\ln\frac{\nu_s}{\nu_n}.\tag{29}$$

Using the value of ν from formula (17), we obtain

$$\delta\dot{\varepsilon}_{ns}=\begin{cases}\dot{\varepsilon}_n(t_0)\dfrac{e^{\Delta/T}-1}{2},\\[2mm] \dot{\varepsilon}_n(t_0)\left(\dfrac{B_n}{M\omega_0}-1\right),\end{cases}\qquad \delta\varepsilon_{ns}=\begin{cases}\dfrac{T}{kv}\ln\dfrac{1+e^{\Delta/T}}{2}, & T_c>T>T_0\\[2mm] \dfrac{T}{kv}\ln\dfrac{B_n}{M\omega_0}, & T<T_0\end{cases}\tag{30}$$

These formulas give relations close to those observed in experiment[4,5].

Let us discuss now the temperature dependence of the elongation $\delta\epsilon_{ns}$ and of the jump of the stress in the case of active deformation $\delta\sigma_{sn}$. As seen from (26) and (30), these quantities differ only by a factor k, which we assume to be independent of the temperature. The temperature dependence of the quantities $v\delta\sigma_{sn}$ and $kv\delta\epsilon_{ns}$ is plotted schematically in Fig. 5, which shows for comparison also the temperature dependence of the superconductor gap $\Delta(T)$. We see that the temperature dependence of the weakening is sensitive to the degree of perfection of the crystal, which is characterized by the length L. Both types of the plots shown in Fig. 5 were observed in experiment[20-23].

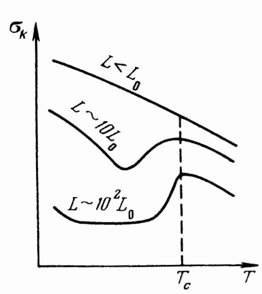

FIG. 2

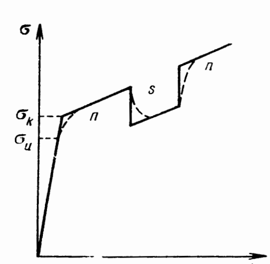

FIG. 3

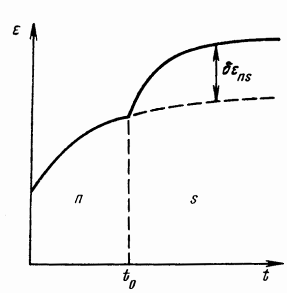

FIG. 4

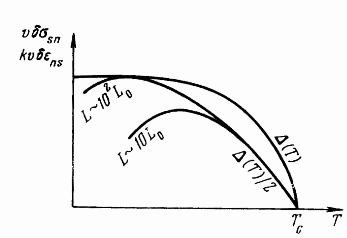

FIG. 5

Stress Relaxation

The total deformation rate $\dot{\epsilon}_u + \epsilon$ becomes equal to zero when the deformation-producing device is stopped instantaneously. The succeeding course of the plastic deformation is due to the elastic stresses present at the instant of stopping, which decrease gradually. The rate of elastic deformation is $\dot{\epsilon}_u = \dot{\sigma}/G$ (G is the shear modulus), and the equation describing the relaxation process can therefore be written in the form

$$\dot{\sigma} + G\dot{\varepsilon} = 0.$$

Substituting here the value of the rate of plastic deformation $\dot{\epsilon}$ from (20), we obtain

$$\dot{\sigma} = -G\varepsilon_0 \nu \exp\left\{-\frac{U - v(\sigma - \sigma_i)}{T}\right\}.$$

Integration of this equation yields the law of the decrease of pressure with time

$$\sigma(t) = \sigma_0 - \frac{T}{v}\ln(\beta\nu t + 1), \quad \dot{\sigma}(t) = -\frac{T\beta\nu}{v(\beta\nu t + 1)}, \quad \beta = \frac{G\varepsilon_0 v}{T}\exp\left\{-\frac{U - v(\sigma_0 - \sigma_{i0})}{T}\right\}, \tag{31}$$

where σ_0 and σ_{i0} are respectively the values of the external and internal stresses at the instant $t = 0$ of the start of relaxation. The degree (depth) of the relaxation is usually characterized by the difference $\Sigma = \sigma_0 - \sigma(t)$ at large values of the time:

$$\Sigma = (T / v)\ln \beta\nu t. \tag{32}$$

We see therefore that the depth of relaxation depends on the state of the metal (on the frequency ν). If one of two structurally identical metals is in the normal state during the time of relaxation, and the other in the superconducting state, then the difference between the depths of relaxation for them is determined, according to (32), by the expression

$$\delta\Sigma_{sn} = \frac{T}{v}\ln\frac{\nu_s}{\nu_n}. \tag{33}$$

This quantity coincides with the jump of the deforming stress $\delta\sigma_{sn}$.

If the transition from the normal state into the superconducting state is produced during the time of the relaxation process (at a certain instant $t = t_0$), then the course of the process changes sharply. In this case the relaxation is described by the formula

$$\sigma_0 - \sigma(t) = \begin{cases} \dfrac{T}{v}\ln(\beta\nu_n t + 1), & t < t_0, \\ \dfrac{T}{v}\ln[\beta\nu_s(t - t_0) + \beta\nu_n t_0 + 1], & t > t_0. \end{cases} \tag{34}$$

Calculating the jump of the relaxation rate at the instant of the ns transition, we obtain

$$\delta\dot{\sigma}_{ns} = \dot{\sigma}_n(t_0)\left(\frac{\nu_s}{\nu_n} - 1\right) = \begin{cases} \dot{\sigma}_n(t_0)\dfrac{e^{\Delta/T} - 1}{2}, & T_c > T > T_0 \\ \dot{\sigma}_n(t_0)\left(\dfrac{B_n}{M\omega_0} - 1\right), & T < T_0. \end{cases} \tag{35}$$

We see therefore that with increasing damping of the relaxation process, i.e., the jump in the relaxation rate decreases with decreasing $\dot{\sigma}(t)$.

The singularities obtained here for the relaxation process are in good agreement with the experimental data[6-8].

Weakening of a Metal on Going from the Normal to the Intermediate or Mixed State

Let us consider a metal in the mixed or intermediate state with a normal-phase concentration c_n. If the dimensions of the normal and superconducting regions exceed the characteristic distances L between the barriers, then the rate of plastic deformation of such a metal will be determined by formula (20), in which one must put

$$\nu = \bar{\nu} = \nu_s + c_n(\nu_n - \nu_s). \tag{36}$$

This relation permits an easy generalization of the results obtained for the ns transition to the case of the ni transition (by ni transition we mean the transition from the normal to the intermediate or mixed state).

For the jump of the deforming stress in the in transition we obtain

$$\delta\sigma_{in} = \frac{T}{v}\ln\left[\frac{\nu_s}{\nu_n} + c_n\left(1 - \frac{\nu_s}{\nu_n}\right)\right]. \tag{37}$$

For the jump of the creep rate and elongation of the sample in the ni transition we have

$$\delta\dot{\varepsilon}_{ni} = \dot{\varepsilon}_n(t_0)(1 - c_n)\left(\frac{\nu_s}{\nu_n} - 1\right), \quad \delta\varepsilon_{ni} = \frac{T}{kv}\ln\left[\frac{\nu_s}{\nu_n} + c_n\left(1 - \frac{\nu_s}{\nu_n}\right)\right]. \tag{38}$$

Finally, the jump of the stress relaxation rate at the instant of the ni transition is determined by the expression

$$\delta\dot{\sigma}_{ni} = \dot{\sigma}_n(t_0)(1 - c_n)(\nu_s / \nu_n - 1). \tag{39}$$

We note that the dependence of the weakening effect in the ni transition on the external magnetic field is determined by the field dependence of the normal-phase concentration c_n.

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263