being increased by the heating contribution as T increases.

[8]The shape of the sample signal obtained with the 6-μsec-pulse laser showed some irregularities which could not be accounted for by transient pickup or the electronic circuitry. A complicating feature of these experiments is that the mean free path of thermal phonons in the sapphire substrate was ~1 cm. Ballistic phonon effects (complicated by a possible dimensional resonance in the substrate) may occur. These effects also hinder a calculation of the thermal response time of the system. The thermal diffusion distance for sapphire in 10 μsec is ~ sapphire phonon mean free path and this is comparable to or greater than the dimensions of the substrate. These phonon effects may be important for the superconducting state of the film and the effects discussed in this paper.

[9]A Kapitza thermal boundary resistance which depends on light intensity, of course, is not precluded.

[10]One may also question the assumption, so far implicit, that the observed effect in the normal state is due to ordinary heating.

[11]Within the skin depth the photons are absorbed at a rate $\sim 10^{24}-10^{25}/\mathrm{cm}^2\,\mathrm{sec}$. If each photon were absorbed by an electron whose excited state lifetime was greater than several microseconds the equilibrium carrier concentration would be altered by less than 1% during the 6-μsec laser pulse. A greater change would occur if each photon led directly or indirectly (e.g., by phonon emission) to much more than one excited electron state. (See further discussion in Sec. IV.)

[12]For a free electron in the alternating electric field of the laser the amplitude of the oscillatory velocity would be ~1 cm/sec. The dc critical velocity for the destruction of superconductivity is $\sim 10^5$ cm/sec.

[13]It would also be important to determine whether ballistic phonons in the sapphire substrate (from the heat pulse) are involved in this effect.

[14]J. R. Anderson and A. V. Gold, Phys. Rev. 139, A1459 (1965).

[15]J. R. Schrieffer and D. M. Ginsberg [Phys. Rev. Letters 8, 207 (1962)] have calculated the recombination time for thermally excited quasiparticles in Pb at 1.44 °K to be ~0.04 μsec. See also A. Rothwarf and M. Cohen, Phys. Rev. 130, 1401 (1963). Quasiparticle lifetimes of ~μsecs have been observed in Al [see K. E. Gray, A. R. Long, and C. J. Adkins, Phil. Mag. 20, 273 (1969)].

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[17]Other processes may contribute to the experimental delay time.

PHYSICAL REVIEW B VOLUME 4, NUMBER 7 1 OCTOBER 1971

Dislocation Inertial Effects in the Plasticity of Superconductors*

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(Received 10 May 1971)

It is shown that radiation damping is weak enough so that all dislocation segments in superconductors are underdamped at low enough temperatures. Consequently, a dislocation impinging on a barrier in the superconducting state overshoots its static-equilibrium position, exerting an additional force on the obstacle, thereby increasing the plasticity. The inertial model gives a quantitative account of available observations on the size of the effect and its dependence on temperature, magnetic field, deformation, and purity.

I. INTRODUCTION

In the past several years, striking observations have been made of an increased plasticity of materials entering the superconducting state.[1–14] These measurements, beginning with the work of Pustovalov *et al.*[1] and Kojima and Suzuki,[3] have recently been reviewed by Alers, Buck, and Tittman.[15] As yet there appears to be no satisfactory quantitative or qualitative explanation. We give here an inertial model[16] of dislocation motion in superconductors and show that it can give a quantitative account of the so-far-available data.

The following are facts which must be explained by an adequate theory:

a. Direction. When a superconducting material is switched into the superconducting state, the plasticity is increased. For constant strain-rate tests, the stress required drops. For creep measurements at constant stress, the strain rate increases dramatically (Soldatov *et al.*[7]). For stress-relaxation experiments at constant strain, the stress drops suddenly (Suenaga and Galligan[10,12]).

b. Magnitude. The stress-change effects observed are typically of the order of from 0.1 to 10%. However, effects as large as 53% have been reported.[5] The effects are strong for Pb, weak for Sn, with Tl and In in intermediate positions (Startsev *et al.*[14]).

c. Universality. The effect appears to be universal. Its existence is independent of crystal structure, appearing in Pb (fcc), Nb (bcc), In (fct), Sn (bct), and Tl (cph). It is found in pure, impure, and alloyed Nb-Mo,[8] Pb-In[14] materials, in single crystals and polycrystals, in weak and strong coupling, and type-I and type-II superconductors.

d. Dependence on deformation. The percentage change in stress is largest at the yield stress, and decreases with work hardening. Over a range of stress of a factor of 15, the stress change has been observed to be nearly constant,[5] increasing through a weak maximum (see Fig. 1). Over narrower stress ranges at lower stress levels, a weak increase in the stress change $\Delta\sigma$ is often found in Pb (Alers *et al*,[6] Suenaga and Galligan,[10,12] and Buck *et al.*[9]). Below the yield stress, no effect is seen.[10,12,14]

e. Temperature dependence. For In, a $1-t^2$, where $t=T/T_c$, temperature dependence for $\Delta\sigma$ is found by Alers *et al.*,[6] and Hutchison and Pawlowicz,[11] while for Pb, a steeper drop near T_c is found by Suenaga and Galligan[12] and Pustovalov and Fomenko.[13] Pustovalov and Fomenko find, however, that their results for In agree with those for Pb. Measurements on Pb by Alers *et al.*[6] at 1.42 and 4.2 °K and by Pustovalov *et al.*[5] at 1.8 and 2.6 °K are in agreement with the Suenaga - Galligan data.

f. Magnetic field dependence. In type-I superconductors, no magnetic field dependence is found apart from that required to change the state.[3,6,7] For type-II superconductors, evidence has been found that $\Delta\sigma$ is proportional to the magnetic induction B for Nb and Nb-Mo alloys (Kostorz[8]), and Pb-In alloys (Startsev *et al.*[14]).

g. Strain-rate dependence. The stress drop $\Delta\sigma$ has been found to be independent of strain rate for Pb for strain rates from 1.6×10^{-6} to 7×10^{-3} sec^{-1},[6,14] for Nb[8] and In.[11]

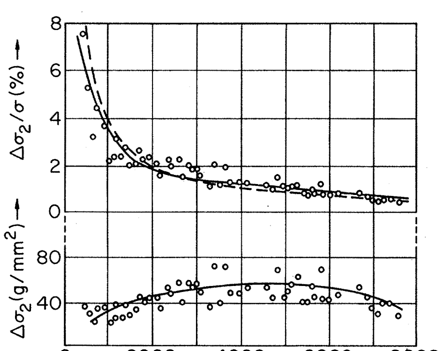

FIG. 1. Flow-stress dependence of the absolute $\Delta\sigma_2$ and the relative $\Delta\sigma_2/\sigma$ value of flow stress increase in polycrystalline lead of 99.9995-at.% purity when the superconductivity destruction takes place. Temperature: 4.2 °K. (After Pustovalov *et al.*, Ref. 5.) The dashed line on the upper curve in the prediction of the inertial-model work-hardening dependence.

h. Impurity and alloying dependence. Impurities and alloying affect the magnitude of the effect; i.e., the effect is not intrinsic. Large increases in $\Delta\sigma$ are found by Kostorz[8] for Nb-Mo alloying. The average creep-rate jump is 32, 143, and 250 for Pb polycrystals, 99.97% single crystals, and 99.992% single crystals, respectively (Soldatov *et al.*[7]). A lower $\Delta\sigma$ was found for purer Pb crystals by Buck *et al.*[9]

i. Reversibility. On reentering the normal state, plastic deformation stops until the stress reaches that level which produces constant strain-rate flow in the normal state (σ_n).[6] In stress-relaxation experiments, the stress drop $\Delta\sigma_r$ decreases with holding time in the normal state. On switching back to the normal state, no further stress change is obtained.[10] No change is found on entering the normal state in creep.[7] A predeformation is required, however, to obtain the strain-rate jump.[14]

j. Crystal orientation. The size of $\Delta\sigma$ is found to depend on crystal orientation at low stresses.[9,12]

The first idea offered[3,5,7] in explanation of the effect was that electron scattering from moving dislocations should provide a viscous drag which would disappear in the superconducting state. The dislocation velocity v would be given by $B_e v = b\sigma$, where B_e is the electronic drag coefficient, b is the Burgers vector, and σ is the applied stress. Then with the usual dislocation strain-rate equation $\dot{\epsilon}=\rho b v$, where $\dot{\epsilon}$ is the strain rate and ρ is the mobile dislocation density, a strain-rate change is predicted at constant stress. However, it has been shown[6] that this model predicts a strain-rate dependence for $\Delta\sigma$ in constant strain-rate measurements which is not observed (in Sec. I*g* above). The argument is not compelling since it is known[17] that for large enough velocities the strain rate is limited not by dislocation velocities but by dislocation generation rates. In fact, it has been suggested[8,14] that the stress change may be associated with a change in the mobile dislocation density. However, the experiments cited (in Sec. I*i*) above are evidence that no changes in mobile dislocation density occur unless the stress is raised above any previous value. This is in agreement with a recent atomistic model for low-temperature creep.[18] Other attempts to save the viscous-drag model rely on appeals to the possibility of a nonlinear dependence of the drag on velocity.[10,12,19,14,20] However, the effect cannot depend only on a viscous drag since it is not an intrinsic effect (Sec. I*h*).

Attempts to explain the effects in terms of structure-sensitive obstacle-dislocation interactions have used[6-9,11] as a basis a rate-theory expression of the form $v=\rho b d\gamma_0\, e^{(V\sigma-U_0)/kT}$, where d is the distance between obstacles, γ_0 is the effective attack frequency,[21] U_0 is the barrier energy height, and V is an activation volume. It is supposed that the barrier energy U_0 somehow changes in the superconducting state. Buck *et al.*[9] suggest that an electrical interaction between impurities and dislocations may change in the transition. Hutchison and Pawlowicz[11] suppose that the activation volume changes. The use of rate theory at helium temperatures may be questioned. In any case, no quantitative theory depending on obstacle drag of this type is available.

It seems plain that neither a purely viscous-drag nor an obstacle model can explain the effects; somehow both are needed. A compromise model has been suggested[22] in which the force on a pinning point is influenced by the thermal vibrations of the dislocation line, which are in turn influenced by the damping of the medium. A model more closely related to that proposed here may be one in the recent work by Kocks and Ashby,[23] who have been considering inertial effects on the motion of kinks and dislocation lines.

II. INERTIAL MODEL

The inertial model we suggest is entirely analogous to a loaded spring in a viscous medium and is easily understood with the help of Fig. 2. A dislocation line moving toward obstacles in position 1 of Fig. 2(a) meets the obstacles with a velocity v_0 at position 2. The static-equilibrium position under an applied stress is position 3. If the viscous damping is larger than a critical value, the dislocation line approaches the static-equilibrium position as in the solid line of Fig. 2(b). If the damping is

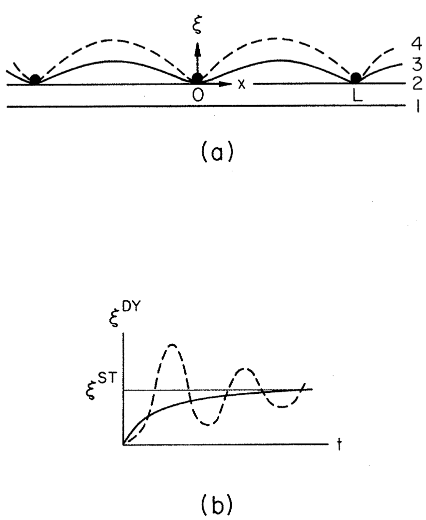

FIG. 2. (a): 1, dislocation line approaches pinning points; 2, dislocation line just touches pinning points; 3, static-equilibrium position of dislocation line; 4, overshoot position of underdamped dislocation. (b) Displacement as a function of time for an underdamped dislocation (solid line).

less than critical, the dislocation line overshoots to position 4, and oscillates about the static-equilibrium line. In position 4, the force exerted by the dislocation line on the obstacle is greater than in the static-equilibirum case. Alternatively, one may say that a smaller stress is needed in the low-damping case to produce the same force on the obstacle. This model in the overdamped case, but not in the underdamped case, has been used by Frost and Ashby[24] and by Klahn *et al.*[25] to compute velocities of dislocations as a function of stress.

To make the model quantitative, we introduce some idealizations. It is assumed that the dislocations are straight when first touching the obstacles and that the string model introduced by Koehler[26] and used extensively in ultrasonic attenuation and internal-friction measurements[27] can be used to describe the motion. In addition, complications arising from changing dislocation tensions with displacements,[28,29] unequal pinning lengths,[30] and finite displacements[24,29] are ignored by assuming constant dislocation tension, equal loop lengths, and weak obstacles. It is assumed that plastic flow proceeds when the depinning force F exceeds a critical value F_c.

The dislocation equation of motion is then given by

$$A\ddot{\xi} + B\dot{\xi} - C\sigma^2\xi/\sigma x^2 = b\sigma\,, \qquad (1)$$

where $A=\rho b^2$ is the dislocation mass per unit length, ρ is the density, B is the viscous-damping constant, and $C\approx Gb^2$ is the dislocation tension where G is the shear modulus. The dislocation displacement ξ is a function of time (derivatives indicated by dots) and position x along the dislocation line. The boundary conditions are $\xi(0,t) = \xi(L,t)$ and the initial conditions are $\xi(x,0)=0$ and $\dot{\xi}(x,0)=v_0$, the initial velocity. The solution ξ can be expressed as the sum of a static displacement ξ^{st} and a dynamic transient displacement ξ^{dy}:

$$\xi = \xi^{\mathrm{st}} + \xi^{\mathrm{dy}}\,, \qquad (2)$$

where

$$\xi^{\mathrm{st}} = b\sigma\, x(L-x)/2C\;. \qquad (3)$$

For sufficiently small damping we have

$$\xi^{\mathrm{dy}} = e^{-\gamma t}\sum_{n=0}^{\infty}\left(B^{(1)}_{2n+1}\cos\omega'_{2n+1}t + B^{(2)}_{2n+1}\sin\omega'_{2n+1}t\right)\sin\left(\frac{(2n+1)\pi x}{L}\right)\,, \qquad (4)$$

where

$$B^{(1)}_{2n+1} = -\frac{4}{\pi^3}\,\frac{b\sigma L^2}{C}\,\frac{1}{(2n+1)^3}\,,\qquad B^{(2)}_{2n+1} = \frac{4v_0}{(2n+1)\pi\omega'_{2n+1}} + \frac{\gamma B^{(1)}_{2n+1}}{\omega'_{2n+1}}\,, \qquad (5)$$

with

$$\omega'_{2n+1}=(2n+1)\omega_0[1-\gamma^2/(2n+1)^2\,\omega_0^2]^{1/2}$$

and $\gamma=B/2A$. The first term of the sum in Eq. (4) gives an adequate representation of the displacement for a simplified discussion. With

$$\omega_0=\frac{\pi(C/A)^{1/2}}{L}\simeq\frac{\pi v_s}{L}\ , \tag{6}$$

where v_s is the shear-wave velocity and $b\sigma L\sim Gb^2$, one obtains $B^{(2)}/B^{(1)}\sim-\pi v_0/v_s+\gamma/\omega_0$. If $v_0\ll v_s$, the $B^{(2)}$ term can be neglected. If $v_0\to v_s$, a relativistic treatment is required in any case. In what follows, we neglect the $B^{(2)}$ term. This will underestimate the displacement, and deviations from the predictions may be examined for evidence of relativistic velocities.

Then we have

$$\xi^{dy}=-(4b\sigma L^2/\pi^3C)e^{-\gamma t}\sin(\pi x/L)\cos\omega' t\ . \tag{7}$$

The depinning force is

$$F=2C\xi_x(1+\xi_x^{\ 2})^{-1/2}\approx 2C\xi_x\ . \tag{8}$$

This has its maximum value when $t=\pi/\omega_0$. Thus the ratio of the dynamic to the static pinning force is

$$F^{dy}/F^{st}=1+(8/\pi^2)e^{-Z}\ , \tag{9}$$

where

$$Z=\gamma\pi/\omega_0=B\pi/2A\omega_0\ . \tag{10}$$

The damping constant B is made up of contributions from radiation damping B_r, electronic damping B_e, and phonon damping B_p:

$$B=B_r+B_e+B_p\ . \tag{11}$$

The phonon damping can be neglected at low temperatures. At $T=0$ in the superconducting state only the radiation damping remains. This has been calculated[31,32] to be

$$B_r=\tfrac{1}{8}\rho b^2\omega \tag{12}$$

and directly measured in ultransonic experiments.[33] For $\omega=\omega_0$, $B_r=\frac{1}{8}A\omega_0$ and $Z_r=\frac{1}{16}\pi$. The critical damping Z_c for an oscillatory transient is given by $Z_c=\pi$, so that all dislocations, regardless of their lengths, are underdamped by a factor of 16 at $T=0$ in superconductors. This accounts for the universality of the observed effects (Sec. I*c*).

Assuming now a background stress σ_0, the force on a pinning point in the superconducting state F_s is given by

$$F_s=b(\sigma_s-\sigma_0)L[1+(8/\pi^2)e^{-Z_s}]\ . \tag{13}$$

The force F_n in the normal state is given by Eq. (13) with the subscript s replaced by n. The relation between σ_s and σ is obtained by equating $F_s=F_n$ (the subscript n is dropped for σ_n). Now, calling

$$\Delta\sigma/\sigma=(\sigma-\sigma_s)/\sigma=(\Delta\sigma/\sigma)_0f(T)=\phi_0f(T)=\phi\ , \tag{14}$$

$$\phi_0=\tfrac{2}{5}(1-\sigma_0/\sigma)(1-e^{-(Z_n-Z_{s0})}), \tag{15}$$

and

$$f(T)=\frac{5}{3}\,\frac{e^{-(Z_s-Z_{s0})}-e^{-(Z_n-Z_{s0})}}{(1-e^{-(Z_n-Z_{s0})})(1+\frac{2}{3}e^{-(Z_s-Z_{s0})})}\ , \tag{16}$$

where Z_{s0} is Z_s at $T=0$, $(8/\pi^2)e^{-Z_{s0}}$ has been taken as approximately $\frac{2}{3}$, and Z_n is assumed to be temperature independent. In the above, the maximum percentage stress change at $T=0$ is given by ϕ_0, while the temperature dependence is carried by $f(T)$ through $Z_s(T)$. The fraction $f(T)$ goes from 1 at $T=0$ to 0 at $T=T_c$.

If the damping in the normal state is such that $(Z_n-Z_{s0})\ll 1$, then Eqs. (15) and (16) simplify to

$$\phi_0=\tfrac{2}{5}(1-\sigma_0/\sigma)(B_{en}\pi/2A\omega_0) \tag{17}$$

and

$$f(T)=(\gamma_n-\gamma_s)/(\gamma_n-\gamma_0)=1-B_{es}(T)/B_e(T_c)\ . \tag{18}$$

If it is assumed that the electronic drag B_e is proportional to the normal electronic density $\rho_n(T)$, then

$$f(T)=1-\rho_n(T)=\rho_s(T)\ . \tag{19}$$

If it is further assumed that the critical depinning force is given approximately by $\alpha U_0/b$, where $\alpha\sim 1$, then

$$b(\sigma-\sigma_0)L=\alpha U_0/b\ . \tag{20}$$

Equations (17) and (20) both contain the two unknowns σ_0 and L. Using Eq. (20) in (17), one obtains

$$\Delta\sigma=\phi\sigma=\alpha U_0B_{en}\rho_s(T)/5b^2Av_s\ . \tag{21}$$

III. DISCUSSION

Equations (14)–(21) summarize the predictions of the inertial model. The stress is predicted to drop by an amount given by Eq. (14) (Sec. I*a*). The maximum effect is given by Eq. (15) (Sec. I*b*). The maximum value is obtained when the damping is large $(Z_n-Z_{s0})\gg 1$ and when the internal stress is negligible. The maximum value is 40%. We interpret measured values larger than this as evidence for dislocation velocities approaching relativistic speeds.

If $(Z_n-Z_{s0})\gg 1$, the temperature dependence as given by Eq. (16) is stronger than so-far-reported values. For the measured values of some tens of percent reported by Pustovalov *et al.*[5] for Pb near the yield stress, the inertial model predicts that the temperature dependence $f(T)$ will be stronger than that for the superconducting electron density $\rho_s(T)$, falling to low values at temperatures below the critical temperature.

For small enough damping in the normal state $(Z_n - Z_r) \ll 1$, the model predicts [Eq. (19)] the temperature dependence of $\rho_s(T)$ for $\Delta\sigma$. In an earlier note,[16] it was found that the data for In by Alers *et al.*[6] and by Hutchison and Pawlowicz[11] are in fairly good agreement with the temperature dependence for $\rho_s(T)$ for weak superconductors according to the BCS[34] theory. The data by Suenaga and Galligan[12] for the strong-coupling superconductor Pb are also in fair agreement with $\rho_s(T)$ determined from the data of Gasparovic and McLean[35] (Sec. I*e*). However, the temperature dependence found by Pustovalov and Fomenko[13] agrees with that of Suenaga and Galligan for both Pb and In.

According to Eq. (17), the percentage effect for small effects ($\ll 40\%$) depends upon the background stress σ_0 and the resonant frequency ω_0, or loop length L. These are unknown. If we assume a simple model for the flow stress [Eq. (20)], we may use Eq. (21) to discuss the dependence on strain hardening (Sec. I*d*). According to the inertial model $\Delta\sigma$ is a constant and ϕ is inversely proportional to the flow stress. In Fig. 1, $\Delta\sigma$ is seen to be approximately, but not quite, constant over a wide range of stress. The model prediction for the percentage change is indicated by the dashed line. The agreement can be regarded as satisfactory, in view of the approximations made. The weakest relation is probably Eq. (20), which ignores many known complications of strain hardening. The model for the change in depinning force is also idealized, but probably gives the change in depinning force more accurately than the absolute magnitude of the depinning force.

For Pb, the constant B_{en}/A has been measured ultrasonically by Hikata and Elbaum.[36] They measured the change in breakaway stress for amplitude-dependent attenuation in the normal and superconducting states. Using a relation for the frequency dependence of the depinning force,[27] they found $B_{en}/A = 2.1\times10^9$ sec^{-1} and $\omega_0 = 3\times10^8$ sec^{-1}. This effect is closely related to the inertial model discussed here. The difference is only that no static force is applied and the breakaway force is provided by a forced oscillation at a much lower stress level in the ultrasonic case. The effect was first observed by Tittman and Bömmel,[37] who interpreted their results in terms of a difference in drag felt by dislocations in the two states.[38] In this sense, this can be regarded as the first observation of the change in effectiveness of pinning points in superconducting transitions.

Using $\Delta\sigma = 40$ g/mm^2,[5] $B_{en}/A = 2.1\times10^9$ sec^{-1}, $v_s = 0.69\times10^5$ cm/sec, $\rho_s(4.2\,^\circ\mathrm{K}) = 0.88$, $b = 3.49\times10^{-8}$ cm, one finds, from Eq. (21), $\alpha U_0 = 0.56$ eV. This appears to be a quite reasonable value. Somewhat smaller values are found from other data at lower stresses. These differences may be interpreted as differences in effective pinning strengths, or possibly a failure of the oversimplified Eq. (20).

Using $\omega_0 = 3\times10^8$ sec^{-1} given by Hikata and Elbaum,[36] one obtains from Eq. (6) a value of $L = 7.2\times10^{-4}$ cm. This is unusually large, but is reasonable for a pure undeformed crystal. If the effective dislocation tension had been overestimated,[28] a somewhat lower value would be derived. Also from the Hikata-Elbaum data, one obtains, from Eq. (10), $Z = 11$, compared to the critical value of π. This means dislocations in this specimen are overdamped in the normal state. In order to obtain a temperature dependence of $\Delta\sigma \propto \rho_s(T)$ as in Eq. (19), it would be necessary for $Z \ll 1$, or L to be about two orders of magnitude smaller. But this is the order of magnitude expected for L in a deformed specimen.

Using $A = \pi\rho b^2$, Hikata and Elbaum derive a value of $B_{en} = 8.6\times10^{-5}$ in cgs units. If we use instead $A = \rho b^2$, we find $B_{en} = 2.7\times10^{-5}$. This is in order-of-magnitude agreement with the calculations of Holstein,[39] Kravchenko,[40] and Brailsford,[41] who predict a temperature-independent electronic drag constant, but is much lower than values calculated by Mason[42] and Huffman and Louat,[43] who predict a temperature-dependent drag constant. It is also much lower than the value determined by Parameswaran and Weertman[44] from their measurements of dislocation velocities under large stresses.

Using $\rho_s(4.2\,^\circ\mathrm{K}) = 0.88$, the electronic damping at helium temperature derived from the Hikata-Elbaum data would be 3.2×10^{-6}. This is still substantially higher than the radiation damping of 5×10^{-7} for such low-frequency loops from Eq. (12). However, for the much smaller loops expected under deformation conditions, the radiation damping will be much greater, and higher temperatures are then needed before the electronic damping becomes more important than the radiation damping.

The magnetic field dependence (Sec. I*f*) finds a natural place in the inertial model since the magnetic induction represents the volume fraction of normal material.[8]

The inertial model as given here is idealized, and the calculations are approximate and incomplete. It is likely that discrepancies which are found can be attributed to approximations in the calculation rather than to a failure of the mechanism. It seems possible that the extra control provided by the ability to vary experimentally the dislocation damping may soon lead to a more complete understanding of the plastic properties of superconductors than of nonsuperconductors. Besides the technical and theoretical interest in ordinary superconductors, the understanding of plastic flow in superconductors is of particular interest for discussions of plastic flow in neutron-star crusts.[45]

*Work supported by the National Science Foundation and the U. S. Atomic Energy Commission under Contract No. AT(11-1)-1198.

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