For real strain rates, reached in the process of removing residual stresses by explosive working of such materials as, for example, St. 3 steel, the dynamic yield stress is 3-4 times greater than the static value. These data, based on experiments with high-speed stretching and compression of rods, are presented in [7].

In addition, the characteristic feature of the dynamic nature of deformation is included by introducing into the analysis the drift of the stressed state of the substance toward the static yield stress.

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RELAXATION OF SUBMICROSECOND PRESSURE PULSES IN A SOLID

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Stress relaxation in dynamical problems of plasticity is described, from the standpoint of dislocation dynamics, by the Sokolovskii–Malvern–Duvall equation [1]:

$$\partial\sigma_{ij}/\partial t - \rho c^2 \partial\varepsilon_{ij}/\partial t = -a\partial\varepsilon^p_{ij}/\partial t \quad (1)$$

which takes into account the effect of velocity on the nature of the wave motion of the deformation. The plastic strain rate tensor $\dot{\varepsilon}^p_{ij}$ is written as the result of the simultaneous gliding in opposite directions of postive and negative dislocations

$$\dot{\varepsilon}^p_{ij} = \sum_{m=1}^{M} \left[{}^+\alpha^{(m)}_{ik} e_{jkl}\, {}^+v^{(m)}_l + {}^-\alpha^{(m)}_{ik} e_{jkl}\, {}^-v^{(m)}_l\right], \quad (2)$$

where the summation is over all the slip planes; ${}^+\alpha^{(m)}_{ik}$ and ${}^-\alpha^{(m)}_{ik}$ are the positive and negative dislocation density tensors.

As a rule, the conditions of deformation at strain rates $\dot{\varepsilon} < 10^3$ ensure, on the average, equality of the positive and negative dislocations, which corresponds to a zero net Burgers vector of the dislocation structure. In the case of pulse or shock loading, however, these conditions may not be fulfilled. In accordance with the definition of the dislocation density tensor in continuum dislocation theory, the latter is written in terms of plastic distortion gradients in the form $\alpha_{ij} = -e_{ikl}\nabla_k w_{lj}$. This means that in the presence of the large displacement gradients realized under high-speed loading the absolute values of the charge dislocation density may also be large. As shown in [2, 3], especially favorable conditions for the appearance of dislocation charges are realized in the contact loading zone.

As is known, charge dislocations are sources of long-range internal stress fields in

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crystals, which may excite collective motions of the dislocation structure. One type of collective motion was examined in [4] with reference to a dislocation wall. In the case of pulse loading the interaction of the stress field and the collective oscillations of the dislocation structure may lead to damping of the loading pulse. This damping may be the cause of additional stress relaxation not taken into account in Eq. (1).

A comparison of the experimental data on stress relaxation associated with high-speed loading and the model represented by Eq. (1) indicates a considerable discrepancy between the stress relaxation rate and the dislocation density. It has been found, for example, that the initial dislocation density, required for the adequate description of the rate of attenuation of the elastic precursor, should be 2-3 orders higher than is observed in the crystals before loading [5], if in Eq. (1) we use the exponential stress dependence of the dislocation velocity proposed by Gilman. The analysis of the damping of the elastic precursor carried out in [6] showed that using the viscous character of the power dependence of dislocation velocity on stress gives realistic values of the dislocation density over the entire region of uniaxial shock loading, except for the contact zone.

In experiments on the damping of the elastic precursor the material is usually loaded by means of high-velocity plate impact. Since very short (<0.1 μsec) pressure pulses cannot be obtained with this method of loading, it is not possible to follow the behavior of the elastic precursor in the contact zone, and the precursor damping curve is extrapolated to the impact plane. At the same time, the latest experiments on loading with very short pressure pulses (30-70 nsec) initiated by lasers [7] or electron beams [8] indicate that in the contact zone the behavior of the compression wave is quite different from what it is in the far zone. It was found, in particular, that in the contact zone there is at first relaxation of the entire wave up to a certain level, followed directly by the separation of an elastic precursor whose damping proceeds at a somewhat different rate. Thus stress relaxation in the elastic precursor is realized from the level to which the entire wave has been able to relax before the elastic precursor was released, and the amplitude of the latter cannot be extrapolated to the impact plane. Thus, the damping of the stress wave is composed of a short wave-damping interval (∿0.5-2 mm) in the contact zone and the subsequent slower damping of the elastic precursor.

A qualitative description of the process of stress relaxation in shock waves from the standpoint of dislocation dynamics was given in [9], where it was shown that a shock wave moving through a crystal contains a charge dislocation surface. In its absence the crystal lattice behind the front would be subjected to very large compressive strains. Under certain conditions the shear stresses acting in planes the normals to which do not coincide with the direction of wave propagation might exceed the theoretical strength of the solid. If there is a moving dislocation surface at the shock front, then behind the front there will be stress relaxation to the hydrostatic compression level.

As shown in [10], longitudinal and shear waves, whose sources are dislocations on the Smith surface, propagate into the region behind the front. These waves transmit the energy of the particles in the wave front into the region behind the front, creating rapidly fluctuating stresses that interact with the dislocation structure of the crystal. The motion of dislocations in random stress fields is characterized by a marked velocity distribution [11, 12], which is indirectly confirmed by measurements of the velocity of the free surface of the target in shock experiments [13].

Thus, on the basis of the above, it is possible to construct the following qualitative picture of stress relaxation in the high-speed loading contact zone: the high charge dislocation densities generated in the wave front create rapidly fluctuating stress fields which interact with the collective oscillations of the dislocation structure, leading ultimately to the dissipation of the energy in the wave. As shown below, the velocity distribution of the dislocations is important in connection with this type of damping.

For the quantitative description of the process of stress pulse relaxation in the contact zone, we will use the known equations of continuum dislocation theory relating the density and charge dislocation flux with the stresses and displacements in the medium [14]:

$$\rho \partial u_k/\partial t = \partial \sigma_{ik}/\partial x_i; \tag{3}$$

$$\partial u_k/\partial x_i = \partial w_{ik}/\partial t - j_{ik}, \tag{4}$$

where u is the particle velocity vector; j_{ik} are the components of the dislocation flux density tensor. A very general relation between j_{ik} and the distortion tensor w_{ik}, taking into account both the spatial and the time dispersion of the waves in a medium with dislocations, was established in [15] with the aid of the so-called dislocation conductivity tensor σ_{iklm}:

$$j_{ik}(x,t) = \int_{-\infty}^{\infty} dx' \int_{-\infty}^{\infty} dt' \sigma_{iklm}(x - x', t - t') w_{lm}(x't').$$

The form of the tensor σ_{iklm} depends on the length distribution of the dislocation segments and their orientation in space. An analysis made in [15] for the particular case of small dislocation oscillations showed that even in the simplest case it is not possible to represent the dislocation conductivity tensor in the form corresponding to a medium with two renormalized Lamé constants, since the contribution of the motion of the dislocations to the different modes of oscillation of the medium is not the same. In this connection it is more constructive to consider the stress field interaction process only in the slip planes of the dislocations, if it is borne in mind that, physically, the dislocations move in precisely this way. In this case the contributions to the total dislocation flux density tensor from the various slip systems can be summed, in the same way as for the total positive and negative dislocation density in expression (2). Differentiating (4) with respect to time and (3) with respect to the coordinate and assuming λ_{iklm} to be constant, we obtain

$$\nabla_i \nabla_l \sigma_{ik} - \rho \lambda_{iklm} \frac{\partial^2 \sigma_{im}}{\partial t^2} + \rho \sum_{m=1}^{M} \frac{\partial j_{lk}^{(m)}}{\partial t} = 0. \quad (5)$$

It is more convenient to proceed using Fourier notation, on the assumption that all the quantities in Eq. (5) vary according to the law exp $(ikx - i\omega t)$. Then Eq. (5) becomes

$$k_i k_l \sigma_{ik} - \rho\omega^2 \lambda_{iklm} \sigma_{im} - i\rho\omega \sum_{m=1}^{M} j_{lk}^{(m)} = 0.$$

In the last term the dislocation flux density tensor, with allowance for the dislocation velocity distribution, can be written in terms of the velocity distribution function [16, 17]:

$$j_{lk}^{(m)} = e_{mli} \tau_l b_k \int_{-\infty}^{\infty} v_m^{(m)} f_{\tau b}^{(m)}(x, v, t)\, dv. \quad (6)$$

Here, $f_{\tau b}^{(m)}$ is a scalar representing the distribution function of the dislocations with respect to the direction of the tangent to the dislocation line τ and the Burgers vector b. In order to find it, we can use the relaxational form of the kinetic equation

$$\frac{\partial f_{\tau b}^{(m)}}{\partial t} + v_i \frac{\partial f_{\tau b}^{(m)}}{\partial x_i} + e_{ikl} \tau_l b_k \sigma_{im} m_{mn}^{-1} \frac{\partial f_{\tau b}^{(m)}}{\partial v_m} = -\frac{f^{(0)} - f_{\tau b}^{(m)}}{\tau_p}, \quad (7)$$

where m_{mn} is the dislocation effective mass tensor; $f^{(0)}$ is the equilibrium velocity distribution function. In order to simplify the further discussion and notation, we will everywhere assume that w_{ik} is the component of the distortion tensor corresponding to the slip plane, and σ_{ik} the corresponding component of the shear stress. Then the components of the distribution function should correspond to the direction of the tangent to the dislocation line τ_i and the Burgers vector component b_k, and the elastic compliance modulus is equal to μ^{-1}, the reciprocal of the shear modulus. In what follows the indices corresponding to the slip plane have been omitted.

We also assume that the distribution function can be represented in the form

$$f = f^{(0)} + f^{(1)},$$

where

$$f^{(1)} \ll f^{(0)}.$$

Then from Eq. (7), in Fourier components, we obtain

$$i\omega f^{(1)} + ikvf^{(1)} + \frac{1}{\tau_p} f^{(1)} = -\sigma b m^{-1} \frac{\partial f^{(0)}}{\partial v},$$

whence

$$f^{(1)} = i\sigma b m^{-1} \frac{\partial f^{(0)}}{\partial v} \left(\omega - kv - \frac{i}{\tau_p}\right)^{-1}.$$

The dislocation flux density tensor (6) can, in its turn, be defined as

$$j = b \int_{-\infty}^{\infty} v f^{(1)} dv,$$

and the dispersion equation for the m-th slip system takes the form

$$k^2 - \frac{\rho}{\mu}\omega^2 + \frac{b^2}{m}\rho\omega \int_{-\infty}^{\infty} \frac{\partial f^{(0)}}{\partial v}\left(\omega - kv - \frac{i}{\tau_p}\right)^{-1} dv = 0. \quad (8)$$

For what follows it is necessary to specify the form of the equilibrium distribution function of the moving dislocations $f^{(0)}$. Along with the common features in the behavior of dislocations and a charged particle gas, in the case of dislocations there are important differences associated with the fact that the dislocations are in another medium. The dislocation kinetics are therefore determined not only by the interaction of the dislocations with each other but also by their interaction with the medium. The latter interaction leads, firstly, to the appearance of dislocation retardation forces due to phonon scattering, phonon viscosity, flutter effect, etc., and, secondly, to the velocity distribution associated with the scatter with respect to the height of the energy barriers in the path of the moving dislocations and, moreover, thermal fluctuations. The equilibrium dislocation velocity distribution function is derived in the appendix. In the unidimensional case it has the following form:

$$f^{(0)} = \left(\frac{2B}{\pi D_2 m}\right)^{\frac{1}{2}} n \exp\left[-\frac{2B}{D_2 m}(v - \bar{v})^2\right]. \quad (9)$$

Here $D_2 = \langle \Delta v \Delta v \rangle / \Delta t$ is the coefficient of diffusion in velocity space, which determines the velocity spread; B is the viscous damping constant of the dislocations. Then Eq. (8) takes the form

$$1 - \frac{\omega^2}{k^2 c_t^2} + \frac{c_p^2}{c_t^2}\frac{\omega}{v_0^3}\int_{-\infty}^{\infty} \frac{v(v-\bar{v})\exp\left[-\frac{(v-\bar{v})^2}{v_0^2}\right]}{\omega - kv - \frac{i}{\tau_p}} dv - \frac{c_p^2}{c_t^2} = 0, \quad (10)$$

where we have introduced the following notation: $c_t^2 = \mu/\rho$ is the shear wave velocity in the crystal; $c_p^2 = \omega_0/k$ is the phase velocity of the oscillations: $\omega_0 = \mu b^2 n/m$ is the natural frequency of the collective oscillations of the dislocation structure; $v_0 = mD_2/2B$ is the mean diffusion velocity of the dislocations. In what follows it is convenient to represent the integral in the form of a sum of two integrals, each of which is the average of the quantities $v^2(\omega - kv - i/\tau_p)^{-1}$ and $v(\omega - kv - i\tau_p)^{-1}$ with respect to equilibrium distribution (9):

$$1 - \frac{\omega^2}{k^2 c_t^2} + \frac{2c_p^2}{c_t^2}\frac{\omega}{iv_0^2}\left[\langle y_2\rangle - \bar{v}\langle y_3\rangle\right] - \frac{c_p^2}{c_t^2} = 0,$$

where

$$\langle y_2\rangle = \frac{1}{\sqrt{\pi}\, k v_0}\int_{-\infty}^{\infty} \frac{v^2 \exp\left[-\frac{(v-\bar{v})^2}{v_0^2}\right]}{\omega - kv - \frac{i}{\tau_p}} dv;$$

$$\langle y_3\rangle = \frac{1}{\sqrt{\pi}\, k v_0}\int_{-\infty}^{\infty} \frac{v \exp\left[-\frac{(v-\bar{v})^2}{v_0^2}\right]}{\omega - kv - \frac{i}{\tau_p}} dv.$$

After averaging we obtain

$$\langle y_2\rangle = \frac{\omega^2}{ik^2}\left[\frac{\sqrt{\pi}}{kv_0} e^{-\left(\frac{\omega}{kv_0} - \frac{\bar{v}}{v_0}\right)^2}\right] + \frac{i}{\omega}\left[1 + \frac{1}{2}\frac{1}{\left(\frac{\omega}{kv_0} - \frac{v}{v_0}\right)}\right], \quad \langle y_3\rangle = 0. \quad (11)$$

With allowance for (11), the dispersion equation takes the form

$$1-\frac{\omega^2}{k^2c_t^2}+2i\sqrt{\pi}\left(\frac{\omega}{kv_0}\right)^3 e^{-\left(\frac{\omega}{kv_0}-\frac{\bar{v}}{v_0}\right)^2}+\frac{c_p^2}{c_t^2}=0.$$

We represent the frequency of the oscillations in the form of a sum of a real and an imaginary part:

$$\omega=\omega_0(1-i\delta).$$

Then in the approximation $\omega_0/k \gg v_0$ for the logarithmic decrement we obtain

$$\delta=\sqrt{\pi}\left(\frac{\omega_0}{kv_0}\right)^3\left\{\exp\left[-\left(\frac{\omega_0}{kv_0}-\frac{\bar{v}}{v_0}\right)^2\right]\right\}\frac{c_p}{c_t}. \tag{12}$$

Here we have neglected the term $(c_t^2/c_p^2)(1-\delta^2)$ in view of the fact that $c_t^2/c_p^2 \ll 1$. This is all the more justified at high values of the decrement when $\delta \to 1$.

Let us consider the two extreme situations corresponding to the absence of viscous damping ($B = 0$) and very high values of the constant ($B = \infty$). In the first case the term in the exponent $\bar{v}/v_0 = (\sigma b/B)(2B/D_2m)^{1/2} \to \infty$, whereas the term $\omega_0/kv_0 = (\omega_0/k)(2B/D_2m)^{1/2} \to 0$. As a result $\delta \to 0$, i.e., there is no damping, as was to be expected. In the second case $\bar{v}/v \to 0$ and $\omega_0/kv_0 \to \infty$, so that the damping is again equal to zero, the absence of damping in this case being due to the immobility of the dislocations. As the diffusion velocity v_0 of the dislocations increases, the decrement falls, which indicates that the dislocation velocity spread leads to a decrease in the damping of the pulse.

With the aid of the expressions obtained we can estimate the viscous damping constant B on the assumption that there is no multiplication of dislocations. The relation between B and the diffusion coefficient in velocity space can be found from the expression [11]

$$D_{2v}=(b^2/B^2)\mu^2b^2n.$$

In this expression, in accordance with [11], D_{2v} is to be understood as the velocity dispersion of dislocations in viscous motion in the random stress field of the crystal. To obtain the coefficient of diffusion in velocity space D_2 it is necessary to average the dispersion D_{2v} with respect to the interaction correlation time. Obviously, in the case of dislocations, as the upper bound of the correlation time we can take the time at which the dislocation reaches the steady-state velocity, which can be determined from the equation of motion of the dislocation

$$m\dot{v}=\sigma b-Bv \tag{13}$$

for known external stress and damping constant. From the equation there follows

$$v=(\sigma b/B)(1-e^{-Bt/m}),$$

whence

$$t_{cor}=m/B.$$

Then the diffusion coefficient D_2 takes the form

$$D_2=(b^4/mB)\mu^2n.$$

From this relation and the definition of v_0 introduced in expression (10) we can define the damping constant as

$$B=\sqrt{2}\sqrt{n}b^2\mu/v_0. \tag{14}$$

The mean diffusion velocity v_0 can be determined with the aid of the expressions previously obtained from experiments on the damping of submicrosecond pressure pulses in the contact zone [7]. The data needed for the calculations are given in Table 1. To these values, in accordance with definitions (10), there corresponds a natural frequency of the collective oscillations of the dislocation structure $\omega_0 = 10^{-10}\,\mathrm{sec}^{-1}$. To the length of the pressure pulse front $\tau_{fr} = 10^{-8}$ sec there corresponds a wavelength of the fundamental harmonic $\lambda = 3\cdot10^{-4}$ cm or a value of the wave vector $k = 2\cdot10^4$ cm, whence the phase velocity $c_p = (\omega_0/k)^{1/2} = 5\cdot10^5$ cm/sec. From expression (12) we can find the ratio c_f/v_0, if the logarithmic decrement δ is known. The latter can be determined from the pressure pulse attenuation curves for aluminum [7]. The pulse amplitude decreases as $\sigma = \sigma_0 \exp[-(\omega_0\delta)t]$.

TABLE 1

Dislocation density	$10^8 cm^{-2}$
Length of loading pulse front	10^{-8} sec
Value of Burgers vector (aluminum)	$2{,}86 \cdot 10^{-8}$ cm
Effective mass of dislocation	$1{,}7 \cdot 10^{-16}$
Shear modulus	25,6 GPa
Mass density of material	2,7 g/cm^3

To the attenuation curve of the pulse with initial amplitude 0.15 GPa in [7] there corresponds the value $\delta = 7 \cdot 10^{-4}$. Then, in accordance with (12), $c_p/v_0 = 3.6$, whence the mean diffusion velocity of the dislocations $v_0 = 1.43 \cdot 10^5$ cm/sec. Substituting this value in expression (14), we obtain $B = 2 \cdot 10^{-5}$ P, which is in good agreement with the published data for aluminum [18]. The estimates show that the mean diffusion velocity of the dislocations in a dynamic compression pulse may be considerable, up to half as much as the velocity of the transverse oscillations.

In the dynamic compression wave the dislocations may have a velocity both less than and greater than the phase velocity of the oscillations. It is known that in an ensemble of particles distributed according to a law of type (9) the number of particles with a velocity less than the given velocity is greater than the number of particles with a velocity greater than the given velocity. Therefore the number of dislocations entrained by the wave exceeds the number of dislocations transferring momentum to the wave. As a result we get damping of the wave, which has been called Landau damping.

Thus, we have shown that the damping of submicrosecond pressure pulses in the high-speed loading contact zone has a collective character similar to that of plasma oscillations.

In conclusion, the author wishes to thank V. I. Vladimirov for useful discussions.

APPENDIX

To find the quasiequilibrium distribution function we use the kinetic equation for a single-particle velocity distribution function with interaction operator in the Fokker–Planck form [16]:

$$\frac{\partial f}{\partial t} - \mathrm{div}_x (vf) + \mathrm{div}_v (\dot{v} f) = -\frac{\partial}{\partial v}(D_1 f) + \frac{1}{2}\frac{\partial^2}{\partial v^2}(D_2 f) \ldots,$$

where D_1 and D_2 are the coefficients of dynamic friction and diffusion in velocity space respectively. In developed form the translational part of the equation includes the term $f\partial\dot{v}/\partial v$; this depends on the derivative of the acceleration with respect to velocity, which can be determined from the equation of motion of the dislocation (13).

Neglecting, in the equilibrium case, the forces of polarization and fluctuation damping as compared with the viscous forces, we obtain the equation

$$\frac{\partial^2 f}{\partial v^2} - 2\frac{\partial f}{\partial v}\left(\frac{\sigma b}{D_2 m} - \frac{Bv}{D_2 m}\right) + \frac{2B}{D_2 m} f = 0,$$

whence, with allowance for normalization, we finally arrive at

$$f^{(0)} = \left(\frac{2B}{\pi D_2 m}\right)^{\frac{1}{2}} n \exp\left[-\frac{2B}{D_2 m}\left(v - \frac{\sigma b}{B}\right)^2\right].$$

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ENERGY LOSS BY PLASTIC DEFORMATION IN RADIAL COMPRESSION OF A CYLINDRICAL SHELL

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INTRODUCTION

During radial compression of a cylindrical liner at a velocity $\lesssim 10^3$ m/sec, its motion differs from that computed by the equations of ideal fluid hydrodynamics. Agreement can be achieved within the limits of measurement error between experimental data and computation if the energy loss by deformation is taken into account [1]. As in this paper, the behavior of a liner fabricated from a homogeneous and isotropic material is considered in [1]. Its length is assumed constant and so large that edge effects can be neglected, and consideration limited to a ring of unit width. In such a formulation the problem of shell wall deformation is equivalent to their uniaxial compression.

The equation of radial axisymmetric motion of a thin liner subjected to external pressure p(t) has the form [1]

$$\rho h\ddot{R} = N/R - p, \tag{1}$$

where N = σh is the circumferential membrane force, σ is the stress in the liner material, h is the shell wall thickness, W = (R_0 − R) is their displacement, R is the radius of the middle

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