$$\frac{Q_I}{Q_s} < 0{,}35e^{\frac{4\zeta_0}{\varepsilon}} + \frac{13}{6}.$$

When $\zeta_0 = -\varepsilon/2$ (extreme case for our press) this gives $Q_I/Q_s < 0.4$ and at twice the velocity ($\zeta_0 = -\varepsilon$) $Q_I/Q_s < 0.06$.

Thus, even for a moderate approach velocity the central unit of the impact device is smaller than that of the static press.

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MECHANISM FOR PLASTIC RELAXATION OF A SOLID IN A SHOCK WAVE

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§1. Model of Phenomenon

We consider plastic relaxation of a solid behind a stationary, plane shock front resulting from above-barrier slip of dislocations. Let the wave be moving in the direction of the x axis at a constant velocity D. We employ a coordinate system moving with the wave and we consider the state of an elementary plane layer of thickness dx which is stationary in this coordinate system. As is usual, we represent the actual dislocation ensemble by four effective slip systems of edge dislocations, the planes of which coincide with the planes of nonzero principal shear stresses (i.e., they make an angle of $\pi/4$ with the planes normal to the coordinate axes). We assume that in any elementary volume and for any slip system, an identical number of dislocations of opposite sign is created per unit time. However, the density of dislocations of opposite sign will not be the same in the elementary layer dx under consideration. Indeed, let the dislocation slip velocity be v. Then (from the assumed stationarity of the wave) the elementary layer dx crosses identical numbers of dislocations of opposite sign per unit time but it crosses them at different velocities: $(D + v/\sqrt{2})$ and $(D - v/\sqrt{2})$, respectively. Therefore, an excess of dislocations moving in the direction of the shock front will be observed in the layer. The relative magnitude of this excess is obviously $(v/D\sqrt{2})$. The effect of an excess of dislocations of one sign is equivalent to the presence in the layer dx of an equivalent Smith wall [1] which is the result of discontinuous relaxation because of a change in the principal strains ε_y and ε_z by an amount $(b\sqrt{2}/l)$, where b is the absolute value of the Burgers vector and l is the distance between dislocations belonging to a single set in the wall. The structure of a layer with a Smith wall is sketched in Fig. 1 with the dislocation density in the wall being exaggerated by several orders of magnitude for clarity. To understand what follows, it is important to emphasize that the Smith wall moves with a velocity D only in the formal sense; in fact, D is the displacement phase velocity of a section in which the dislocation density in the wall has a certain definite value whereas the excess dislocations themselves move with a velocity v. Besides the two sets of dislocations shown in Fig. 1, the wall contains yet another two sets of dislocations which are parallel to the plane of the figure so that the total number of dislocations per unit area of the wall is $4/l$. In the relief region behind the compression wave, the sign of v changes and the direction of the Burgers vector for the excess dislocations and for the Smith

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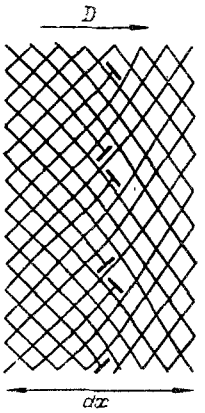

Fig. 1

wall changes sign. The quantity l is connected with the layer thickness dx and the density N of effective dislocations (which is close to the true density of moving dislocations) by the relation

$$ldx = 2\sqrt{2}D/Nv. \tag{1.1}$$

We now replace the continuous change in state parameters of the material in the layer dx by their discontinuous change in the equivalent Smith wall. For computational simplicity, we confine ourselves to consideration of small strains $\varepsilon_i \ll 1$, i.e., shock waves of moderate intensity in which the pressure is much less than the bulk modulus of the material. We denote state parameters ahead of the wall by quantities with the subscript 0 and those behind the wall by quantities without a subscript. We keep in mind that $\varepsilon_y = \varepsilon_z$. Then the conditions for conservation of mass, momentum, and energy at the discontinuity are written in the form

$$\begin{gathered}\rho_0 D = \rho(D-u), \text{ i.e.,}\\ u = D[(\varepsilon_{x0}-\varepsilon_x)+2(\varepsilon_{y0}-\varepsilon_y)],\ \sigma_{x0}-\sigma_x = \rho_0 uD,\\ -\sigma_x(D-u)+\rho_0 D\left[\frac{(D-u)^2}{2}+c_0T\right]+\frac{D-u}{2}(\sigma_x\varepsilon_x+2\sigma_y\varepsilon_y)-\\ -4\sqrt{2}\,\frac{\varepsilon D^2}{lvN}\,\frac{dN}{dx} = -\sigma_{x0}D+\rho_0 D\left(\frac{D^2}{2}+c_0T_0\right)+\frac{D}{2}(\sigma_{x0}\varepsilon_{x0}+2\sigma_{y0}\varepsilon_{y0}).\end{gathered} \tag{1.2}$$

One should add to them Hooke's law for a plane wave

$$\begin{gathered}\sigma_x = \frac{E}{1-2\nu}\left[\frac{(1-\nu)\varepsilon_x+2\nu\varepsilon_y}{1+\nu}-\alpha T\right],\\ \sigma_y = \sigma_z = \frac{E}{1-2\nu}\left[\frac{\nu\varepsilon_x+\varepsilon_y}{1+\nu}-\alpha T\right]\end{gathered} \tag{1.3}$$

and the condition

$$\varepsilon_{y0}-\varepsilon_y = \varepsilon_{z0}-\varepsilon_z = b\sqrt{2}/l, \tag{1.4}$$

where u is the jump in mass velocity at the discontinuity; ρ is density; E is Young's modulus; ν is the Poisson coefficient; c_0 is heat capacity; α is the coefficient of thermal expansion; T is the temperature measured from the state $\sigma_i = \varepsilon_i = 0$; σ_i are the principal stresses; ε is the formation energy for unit length of a dislocation filament. The solution of the system (1.2)-(1.4) is extremely laborious in the general case. For small strains, however, thermal expansion can be neglected by setting $\alpha = 0$ without entailing much loss of accuracy. By solving the system (1.2)-(1.4) in this approximation, we find

$$\rho_0 c_0(T-T_0) = \frac{4b\sqrt{2}}{l}\left(-\tau-\frac{\varepsilon D}{bNv}\,\frac{dN}{dx}\right); \tag{1.5a}$$

$$\sigma_{x0} - \sigma_x = \frac{2\sqrt{2}E}{1-2\nu}\frac{b}{l}\frac{2\nu(1-\nu)-(1-2\nu)(\rho_0 D^2/E)}{1-\nu-(1+\nu)(1-2\nu)(\rho_0 D^2/E)}, \tag{1.5b}$$

where τ is the principal shear stress

$$\tau = \frac{E}{2(1+\nu)}(\varepsilon_x - \varepsilon_y)$$

(in a compression wave $dN/dx < 0$ and $\tau < 0$). Equation (1.5a) is valid as long as the density of moving dislocations in the wave increases; if it falls, the term containing (dN/dx) in this equation should be neglected.

§2. Law of Motion for Dislocations

It is natural to consider the increase in normal stress $(\sigma_{x0} - \sigma_x)$ in the layer dx as the sum of projections on the x axis of forces applied to the Smith wall, i.e., to the excess dislocations of dominant sign. The velocity component of the latter in the x direction is $v/\sqrt{2}$ and the power expended by an external source in maintaining the dislocations is $(\sigma_{x0} - \sigma_x)v/\sqrt{2}$. Since all moving dislocations are equivalent (it is impossible to indicate which of them make up the Smith wall), this power is distributed uniformly among all Ndx moving dislocations. Equating the quantity $(\sigma_{x0} - \sigma_x)v/\sqrt{2}$ to the total power needed to maintain motion of the dislocations at the velocity v, some information can be obtained about the specific nature of the laws of motion for dislocations in a shock wave.

We use the law of viscous above-barrier slip for the dislocations; it is valid for stresses very much greater than the static yield point of the material

$$\tau b = Bv, \tag{2.1}$$

where B is the coefficient of viscosity. The force acting on a dislocation is τb [2] and the power dissipated by a moving dislocation is τbv. Then from the equation

$$(\sigma_{x0} - \sigma_x)v/\sqrt{2} = \tau bvNdx$$

and from Eqs. (1.1), (1.5b), and (2.1) there follows

$$B = \frac{bE}{2\sqrt{2}D(1-2\nu)}\frac{2\nu(1-\nu)-(1-2\nu)(\rho_0 D^2/E)}{1-\nu-(1+\nu)(1-2\nu)(\rho_0 D^2/E)}. \tag{2.2}$$

Thus, the quantity B takes on a specific value in this case which is typical of motion of a dislocation in a stationary shock wave and which does not agree with the phonon viscosity B_p that appears during the motion of a dislocation in a quasistatically loaded crystal. The reason for this difference is that a dislocation moving in a shock wave experiences acoustic resistance that is considerably greater than the viscous frictional resistance in a phonon gas.

In fact, one can rewrite Eq. (2.2) in the form

$$B = \frac{\rho b c^2}{2\sqrt{2}D}\frac{1+\nu}{1-\nu}\frac{2\nu(1-\nu)-(1-2\nu)(\rho_0 D^2/E)}{1-\nu-(1+\nu)(1-2\nu)(\rho_0 D^2/E)},$$

where c is the velocity of longitudinal acoustic waves. Ordinarily, c and D are close in value and one can therefore set $B = m\rho bc$, where m is a dimensionless coefficient with a value of the order of 10^{-1} depending on the specific conditions of the problem. We then have from Eq. (2.1)

$$\tau = m\rho vc.$$

The last equation is similar to the equation determining the pressure in a weak shock wave excited by a piston moving with a constant velocity $v \ll c$. Actually, in this problem, the dislocations act like pistons each of which advances an odd atomic half-plane of the Smith wall in the crystal at a velocity $v/\sqrt{2}$ (the total effect of all $D\sqrt{2}/v$ dislocations occurring at a single dislocation of the Smith wall leads to the fact that the odd atomic half-plane advances in the crystal with a velocity D).

§3. Wave Profile Equations

By substitution of Eqs. (1.1), (2.1), and (2.2) into Eq. (1.5b) we obtain

$$d\sigma_x/dx = -2b\sqrt{2}N\tau \tag{3.1}$$

(in a compression wave, $\tau < 0$ and the compressive negative stress increases in absolute value in the direction of negative x while the wave moves in the direction of positive x). Having found the increase $(\varepsilon_{x0} - \varepsilon_x)$ from Eqs. (1.2)–(1.4), we can obtain

$$\frac{d\tau}{dx} = -\frac{1-2\nu}{\sqrt{2}(1+\nu)}\frac{1-3(1-2\nu)(\rho D^2/E)}{2\nu(1-\nu)-(1-2\nu)(\rho D^2/E)}, \tag{3.2}$$

and, finally, we have from Eq. (1.5a)

$$\rho c_0 dT/dx = -2Nb^2\tau^2/BD - 2\varepsilon dN/dx.$$

The quantity bN has the sense of an inverse width for the zone of plastic relaxation.

In the derivation of general equations for the profile of nonstationary elastoplastic waves, the arbitrary amplitude of the deformation is ordinarily represented as the sum of elastic and plastic components and the Orovan relation $\varepsilon^P \approx (\frac{1}{2})bN\nu$ [3, 4] is assumed for the plastic deformations. Only elastic deformations were considered in the derivation of Eqs. (3 1) and (3.2), and the Orovan relation was not used. This provides an opportunity for indirect evaluation of the accuracy of the approximation under which Eqs. (3.1) and (3.2) were obtained. Use of the Orovan relation leads to the expression

$$D = \sqrt{E/2\rho(1+\nu)(1-2\nu)}.$$

The quantity D is a constant in accordance with the nature of the approximation used (neglect of terms of the order of ε_1) and is close to the velocity of sound. Some deviation of D from the true value is unimportant if one considers the low accuracy with which the quantities N and v are known at the present time.

§4. Possibility of Comparison with Experiment

The two unknown functions N and v appear in the general equation for the profile of elastoplastic waves. The function v is defined in Eqs. (3.1) and (3.2). Equations (3.1) and (3.2) can be considered as particular cases of the general equations for weak stationary waves. Indeed, from a comparison of Eqs. (3.1) and (3.2) with the corresponding simplification of the general equations in [3] [for the slip law (2.1)], it turns out that only the values of the coefficients for the ratio (N/B) are different. The value of this coefficient in Eqs. (3.1) and (3.2) is larger by approxiamtely the same factor by which B is larger than B_p; therefore, calculations of the profile of weak waves from the equations in [3] with phonon viscosity leads to the same results as calculations based on Eqs. (3.1) and (3.2) with a somewhat changed N.

Estimates of the value of B for steel yield values of the order of 10^{-2} p, which is somewhat greater than B_p but agrees with the estimates in [5, 6]. Direct experimental verification of Eq. (2.2) is hardly possible at the present time since the unknown function N remains in these equations.

The problem of the choice of N was discussed in [3, 4, 7] and a number of other papers; it is far from resolution at the present time and therefore the most interesting method of using Eqs. (3.1) and (3.2) is for numerical analysis of experimental data for the purpose of determining N.

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APPLYING THE KINETIC THEORY OF STRENGTH TO DETERMINING LIFE OF LIGHT ALLOYS AFTER PRELIMINARY PLASTIC DEFORMATION

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Residual plastic strains can occur in light-alloy structural elements as a result of the action of a complex of service loads (especially at points of stress concentration) and also in connection with certain fabrication processes involving bending, stretching, various types of cold forging, etc.

The temperature conditions and the amount of plastic deformation determine the previous history of the material and may have an important influence on the resistance of the alloy to static and variable loads. It has been experimentally established that the nature of the effect of plastic deformation depends on many factors: the composition of the alloy, the conditions of preliminary deformation, and the load-temperature conditions of the subsequent tests. Depending on the circumstances, the strength properties of the material may increase, decrease, or remain unchanged.

The extensive use of plastically deformed materials has led to many attempts to determine experimentally the optimum plastic strain. However, the purely experimental approach is very laborious and does not always give reliable results for the favorable degree of plastic deformation.

We have investigated the effect of the degree of preliminary plastic tensile deformation on the long-time strength of Duralumin at various temperatures and stresses and have estimated the possibility of determining the long-time strength under these conditions on the basis of the kinetic theory of strength.

As the material for investigation we used sheets of commercial hardened and naturally aged D16AT alloy 3 mm thick. All the sheets were taken from the same melt. The specimens were cut parallel to the direction of rolling. The mechanical properties of the material in the starting state (as delivered) were as follows: ultimate strength σ_u = 48.6 kg/mm^2, yield point $\sigma_{0.2}$ = 36.4 kg/mm^2, relative elongation δ = 14.3%.

The mechanical characteristics were determined and the preliminary plastic deformation (2, 4, 6, and 8%) was produced at room temperature in a GURM-10 hydraulic tensile testing machine. The specimens subjected to plastic deformation were finished specimens which after cold-hardening had not undergone any further mechanical working. Due to the specially designed recording instrument, the plastic deformation could be carried out with a high degree of accuracy. The spread of the valves at each plastic strain level did not exceed 0.01-0.03%. After preliminary deformation the specimens were tested to destruction in creep at temperatures of 125, 150, 175, and 200°C and stresses of 18, 20, 22, 24, 26, 28, 30, 32, 34, 36, 38, 40, and 42 kg/mm^2.

The tests were conducted on redesigned DST-5000 machines; undeformed specimens were also tested to destruction. Altogether 235 specimens were tested.

The test conditions and the corresponding values of the stress rupture life for various levels of preliminary plastic deformation (PPD) are indicated in Table 1, from which it is possible to judge the complex influence of the preliminary plastic deformation on the long-time strength. Thus, at 125 and 150°C and stresses σ = 34, 36, 38, 40, and 42 kg/mm^2, PPD

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