THE PHYSICAL NATURE OF CREEP IN CRYSTALLINE BODIES. A REVIEW

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1. Introduction. Creep is a special case of plastic deformation in which the external load is kept constant throughout the test. The loading process itself is accompanied by elastic and "instantaneous" plastic deformation. This is followed by the "primary" stage of creep. In this stage the rate of plastic deformation usually decreases (curve a in Fig. 1), although in some cases (in relatively undislocated crystals of silicon and germanium, in polycrystalline iron with dislocations blocked by carbon, in brasses, and other materials) an incubation period is observed (curve b in Fig. 1). During the "secondary" stage (steady-state creep) the creep rate is almost constant. The third "tertiary" stage is usually characterized by an increase in creep rate and ends in fracture.

Like any form of plastic deformation of crystalline material, the atomic mechanism of creep consists in the motion of various lattice defects: point defects (vacancies, interstitial atoms, impurity atoms), linear defects (dislocations), and surface defects (block boundaries, twinning).

In what follows we describe the present state of physical research (mainly in relation to pure metals) into the mechanisms controlling the creep rate under various conditions and the nature of the structural changes taking place during creep. Unfortunately, the results obtained so far are too few for a complete physical theory of creep to be constructed. However, it is already possible to discuss certain of its aspects.

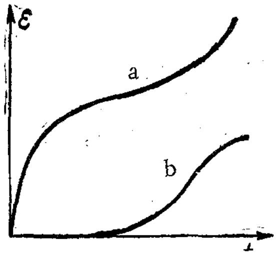

Fig. 1.

2. Motion of point defects (diffusion creep). Migration of point defects occurs in the presence of a chemical potential gradient:

$$\mu = kT \ln \frac{c}{c_0} + p\Omega , \qquad (2.1)$$

where c is the local concentration of defects, c_0 their equilibrium concentration at temperature T, k the Boltzmann constant, p the pressure, and Ω the atomic volume. The mass transfer that determines the plastic deformation is described by he diffusion equation:

$$\frac{\partial c}{\partial t} = D\nabla^2\left(c + c_0 \frac{p\Omega}{kT}\right), \qquad (2.2)$$

where D is the diffusion coefficient of the defects. Of course, the case of vacancies mass transfer (migration of atoms) takes place in a direction opposite to that in which the vacancies move (Fig. 2).

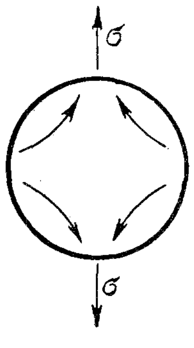

Fig. 2.

Diffusion plasticity can occur both in a nonhomogeneous stress field with grad p ≠ 0 (upward diffusion [1]) and in a homogeneous stress field. In the latter case the defect concentration gradient is ensured by the boundary conditions: difference in chemical potential at different points on the surface of the crystal, or grain in the case of a polycrystal. The stress normal to the surface of the body

$$\sigma_n = \sigma_{ik} n_i n_k$$

(n – normal to surface, recurring indices imply summation) changes the chemical potential by an amount:

$$\Delta\mu = \sigma_n \Omega \qquad (2.3)$$

corresponding to the supersaturation:

$$\Delta c = c_0 \frac{\sigma_n \Omega}{kT}.$$

For a spherical grain of radius R the creep rate is [2, 3]:

$$\dot{\varepsilon}_{ij} = \frac{1}{\eta}\left(\sigma_{ij} - \frac{\delta_{ij}}{3}\sigma_{kk}\right)\left(\eta = \frac{3}{4}\frac{kT}{D\Omega}R^2\right). \qquad (2.4)$$

(δ_{ij} — Kronecker delta).

Taking into account the relaxation of shear stresses in the polycrystal due to the reduced viscosity of the boundary layers introduces a slight change in the numerical coefficient of (2. 4).

Recently, Lifschitz [4] gave a complete solution of the self-consistent problem of diffusion creep in polycrystals which takes into account the internal stresses due to nonconformity of deformation of adjacent grains. For the effective viscosity, in place of the quantity η in Eq. (2. 4), Lifschitz obtained an expression of the type $\eta^\circ + \eta^*$, where η° is determined by the dependence of the bulk viscosity on the shape of the grain, while η^* describes the viscous slip along the grain boundaries. If the latter are sufficiently distorted, the viscosity tensor

$$\eta_{iklm} \approx \overset{\circ}{\eta}_{iklm} = \frac{kT}{4\pi D\Omega V} \int\limits_{V}\int\limits_{V} x_i \xi_k \frac{\partial^2}{\partial x_l \partial \xi_m} |\mathbf{x} - \boldsymbol{\xi}|^{-1} \, d\mathbf{x} d\boldsymbol{\xi}\,, \tag{2. 5}$$

where V is the grain volume. With very fine grains the role of surface diffusion along the grain boundaries becomes more important. If the ratio of the coefficients of surface to volume self-diffusion D_S/D is much greater than the ratio of the grain size to the interatomic distance L/a, then the effective viscosity

$$\eta = \eta^\circ \frac{DL}{D_S a}\,. \tag{2. 6}$$

It is assumed that diffusion creep determines the behavior of fine-grained materials at very high temperatures [5]. The creep of polycrystalline aluminum oxide in the temperature range 1400-1900°C [6, 7] may be quoted as an example; here the relation $\dot{\varepsilon} \approx \sigma D/R^2$, characteristic of diffusion creep, is quite distinctly observed, and the effective viscosity corresponds to the diffusion coefficient for aluminum ions, exceeding by several orders that for oxygen ions. It is possible that the diffusion mechanism of deformation was observed in tensile tests on thin iron wires at temperatures exceeding 1350°C [8]. In these tests grooves, nearly normal to the axis, were formed in the boundary region, and the elongation was linked with the volume of metal expelled from these grooves.

3. Conservative motion (slip) of dislocations. Metallographic investigations show that creep is usually accompanied by the formation of slip lines on the surface. In aluminum single crystals creep due to dislocation slip, in a few clearly visible bands has been observed even at temperatures just below the melting point [9]. When a dislocation moves, the rate of plastic deformation is described by an equation of the type (for details see [10]):

$$\dot{\varepsilon} = Nvb\,, \tag{3. 1}$$

where N is the density of the dislocations, v is their rate of motion, and b is the Burgers vector. The creep curves for germanium can be satisfactorily described by Eq. (3. 1) if we introduce relations $N(\varepsilon)$ and $v(\sigma, T)$ determined from microscopic observations of dislocations [11].

Presumably, at sufficiently low stresses the mobility of the dislocations is determined by the lattice friction, i. e., by forces of the Peierls-Nabarro type (see [12], p. 104), which must be overcome when a dislocation line moves in the slip plane from one position to another with minimum expenditure of energy. The linear sections of a dislocation, lying in favorable directions (in the case of the fcc lattice these are the directions of close packing <110>), are connected by kinks. The tangential displacement of kinks under the influence of stress acting at a rate v_k causes the translational movement of the dislocations at a rate:

$$v = nbv_k\,, \tag{3. 2}$$

where n is the linear density of the kinks. The equilibrium density of the kinks is [13]:

$$n = \frac{(2\pi\rho w_0 kI)^{1/2}}{hE^{1/2}} \exp\frac{-w_0}{kT}\,, \tag{3. 3}$$

where w_0 is the kink energy, ρ the density of the material, E the energy of the dislocation per unit length, and h Planck's constant. According to [14],

$$v_k \approx \frac{C\sigma b^3}{kT}\,, \tag{3. 4}$$

where C is the speed of sound. Consequently, according to (3. 1)-(3. 3), in the extreme case of low stresses dislocation slip leads to a linear relation between creep rate and stress:

$$\dot{\varepsilon} = \frac{N(2\pi\rho w_0)^{1/2} Cb^5}{h(EkT)^{1/2}} \sigma \exp\frac{-w_0}{kR}\,. \tag{3. 5}$$

Evidently, lattice friction controls the creep of covalent crystals [diamond, silicon ([12], p. 505)], and the hexagonal metals zinc (see [15]) and magnesium [16] in the case of slip in non-basal systems. Owing to the high Peierls stresses these systems become active only at high temperatures.

The motion of dislocations may be resisted not only by interaction with the lattice but also by interaction with other dislocations and other crystal defects. For example, in the case of the intersection of mutually attracting dislocations the dislocation reactions take place with the formation of a resultant dislocation favorable with respect to energy, i. e., AB in Fig. 3, which illustrates the interaction between a moving dislocation with Burgers vector b_1 and counter-dislocations; the arrows indicate the direction of the dislocation lines; $b_1 + b_2 + b_3 = 0$. If the dislocation is to advance further, the segment AB must be shortened and the corresponding barrier potential overcome [17]. When mutually repelling dislocations intersect, no reactions take place between them. The moving dislocation may be forced through the wall of the sessile dislocations. In this case the barrier potential U and its dependence on σ are determined by the geometry of the problem. If, instead of making an exact calculation of the elastic fields, we simply compute the energy per unit length (linear tension) of the dislocations, then, in states of stable and unstable equilibrium, the segment punched through is an arc of radius $R = 0.5Gb/\sigma$ ([18], p. 70).

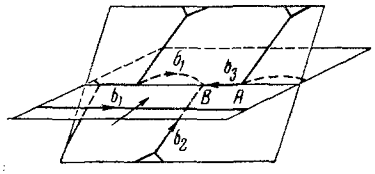

Fig. 3.

The barrier potential is overcome, without activation, at the critical stress

$$\sigma = Gb/l, \tag{3.6}$$

where l is the distance between the wall dislocations.

The slip of screw dislocations in a face-centered lattice is crystallographically possible along two planes of the type {111} passing through the dislocation line. Therefore a screw dislocation, blocked by some obstacle (parallel dislocation, separated particle), can continue its motion by crossing to another plane (cross slip). A favorable dislocation configuration in the face-centered lattice, from the energy point of view, is a stacking fault in the slip plane {111} bounded by partial dislocations (for details see [19], p. 183, and [20]). For a dislocation line to cross from the initial slip plane into a cross-slip plane it is necessary to compress the stacking fault over a certain length at the expense of thermal fluctuation. The compressed section of the dislocation then breaks down into partial dislocations but in a different slip plane; for example, Fig. 4 shows schematically the cross slip of a screw dislocation from the initial plane (1) into plane (2). An attempt to compute the activation energy of this process theoretically has been made by Shoeck and Seeger [21] and Friedel [22].

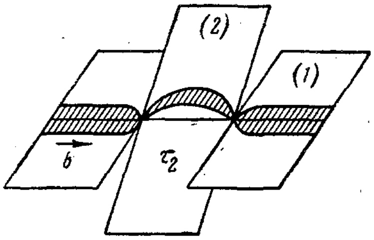

Fig. 4.

If a split screw dislocation has a jog, cross slip will be much easier. The jog in a screw dislocation is dissociated in such a way that one of its partial dislocations already lies in the cross-slip plane (for details see [23]).

It may be assumed that the cross slip of screw dislocations controls creep in polycrystalline aluminum at 275-370°K (activation energy 28 000 cal/mole) [24] and in single crystals of aluminum in pure shear at 600-750°K [25].

4. Non-conservative motion (climb) of dislocations. In contrast to the conservative motion (slip) considered above, non-conservative motion is characterized by the growth or shortening of incomplete atomic planes requiring mass transfer – diffusion of vacancies or interstitial atoms. As a result, the whole dislocation, or parts of it (down to atomic jogs), "climbs," i. e., moves at an angle to the slip plane.

Jogs in dislocations are produced both by intersection with other dislocations whose Burgers vectors do not lie in the same slip plane ([18], p. 200) and by the motion of kinks of different types along a screw dislocation [26].

The activation energy of jog formation on intersection of dislocations depends on the type of jog and the degree of splitting of the dislocations; it is given by $U = \alpha G b^3$, where α may vary from 0.2 to 1 ([18], p. 201). According to Seeger ([19], p. 211), the activation volume γ is then equal to bdl, where l is the distance between the dislocations of the "forest," and d is the width of the split dislocation. The creep rate is

$$\dot{\varepsilon} = \nu_0 bFQ \exp\left[-\frac{U - \gamma(\sigma - \sigma_i)}{kT}\right], \tag{4.1}$$

where ν_0 is the frequency of the thermal vibrations of the atoms (of the order of 10^{12} $\sec^{-1}$), Q is the bulk density of the segments of the dislocations blocked by obstacles, F is the area covered by these segments after overcoming the obstacle, and σ_i is the amplitude of the internal stress field blocking the motion of the dislocations. Assuming that the internal stress is proportional to the strain ($\sigma_i = \theta\varepsilon$) and that the coefficient of the exponent in (4.1) remains constant, after integrating we get the logarithmic relation, frequently observed experimentally,

$$\varepsilon = A \ln(\nu t + 1)$$

$$\left(A = \frac{kT}{\theta\gamma},\ \nu = \frac{1}{A}\nu_0 bFQ \exp\left(-\frac{U-\gamma\sigma}{kT}\right)\right). \tag{4.2}$$

Agreement with the experimental data is obtained if $\gamma \simeq 10^{-19}$ cm^3, which corresponds to the observed values of l.

The creep will be steady if all the quantities entering into (4.1) remain constant. A dependence of the steady-state creep rate on stress and temperature of type (4.1) has been observed experimentally by Zhurkov [27, 28] for various materials at temperatures below 0.5 T_* (melting point) and not too low stresses. For example, Fig. 5 [27] gives creep rate versus stress for polycrystalline aluminum: 1 – 100°C, 2 – 200°C, 3 – 300°C, 4 – 400°C, 5 – 500°C. For uniaxial tension, the activation energy U is of the same order as the sublimation energy (about 8.6 Gb^3 for face-centered metals). Unfortunately, the atomic mechanism that limits the creep rate has not been determined either by Zhurkov nor by many other workers. However, according to Myshlyaev's electron microscope data [29], l is actually close to the distance between the dislocations at subboundaries: l equals 140 A for γ = 3.66 (kcal/mole)(mm^2/kg).

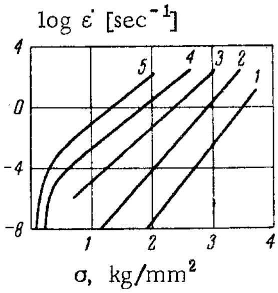

Fig. 5.

A relation of type (4.1) is retained for torsion [30, 31], but in this case the activation energy is much smaller than the energy of sublimation (and even smaller than the energy of self-diffusion). The dependence of the creep rate on the type of stress state was also noted in [32]: the creep rate was 2-3 times greater in compression than in tension.

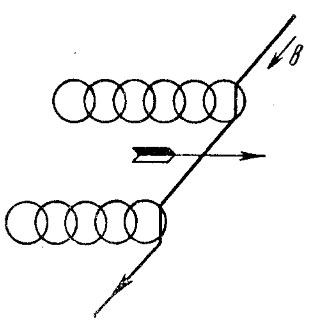

Fig. 6.

When the stress state changes so does the relation between the activity of the various slip systems and, consequently, we get a different dislocation structure and, in particular, a change in the proportion of jogs of various types. It is possible that this also explains the effect of the type of stress state on the activation energy and creep rate.

A number of authors believe that the creep rate is limited by the formation and subsequent motion of jogs resulting from the intersection of dislocations [33, 34]. These may migrate along the dislocation line or climb, leaving behind chains of vacancies or interstitial atoms (see Fig. 6). Correspondingly, if the direction of motion of the dislocation is fixed, we must distinguish jogs of the vacancy and interstitial types. Since the energy of formation of interstitial atoms considerably exceeds the energy of formation of vacancies we may assume that interstitial atoms will only be formed if the motion of the dislocations is rapid. Otherwise the non-conservative motion of an interstitial jog might be maintained by absorption of vacancies reaching the jog from the interior of the crystal or along the dislocation. According to Mott [33], the creep rate is limited by the non-conservative motion of vacancy jogs and is described by an equation of type (4.1), where U is treated as the energy of vacancy formation, and $\gamma = b^2 l$ (l is the distance between jogs). If after the formation of a vacancy the jog does not slip sideways, energy must be spent in removing the vacancy, and the activation energy U will then be composed of the energy of formation of the vacancy and its energy of migration, i.e., it will be equal to the energy of self-diffusion U_d. Feltham [34] has examined in detail various mechanisms of jog motion and has given the spectrum of the corresponding activation energies.

The formation of vacancies in creep is supported by many indirect data. However, it should be pointed out that in plastic deformation in certain circumstances interstitial atoms rather than vacancies are chiefly formed [35].

The above discussion of the effect of the stress state on the type of creep is fully applicable to the case of creep limited by the formation of point defects.

If for some reason dislocation slip is suppressed, the main contribution to creep is made by climb. In this case point defects may be exchanged both between dislocations and between dislocations and crystal (grain) boundaries. At the edges of extra planes the chemical potential varies according to Eq. (2.3), where n is the normal to the extra plane. The motion of point defects toward the edge of the extra plane (or away from it) enables the dislocation to climb at the rate:

$$v_n = \frac{2\pi D}{b \ln(R/r_0)} \frac{\sigma\Omega}{kT}, \tag{4.3}$$

where R is the distance between the source and "sink" of the defects, and r_0 is the radius of the dislocation loop. Substitution of (4.3) in (3.1) gives the creep rate for a crystal with a dislocation density N:

$$\dot{\varepsilon} = \frac{2\pi N\sigma}{\ln(R/r_0)} \frac{D\Omega}{kT} \tag{4.4}$$

Comparison of (4. 4) with (2. 4) shows that even a few dislocations in the interior of a grain ensures a higher creep rate than the pure diffusion (without dislocations) mechanism considered in §2.

To climb the dislocations must have jogs with which vacancies are associated. The activation energy for climb U_c is equal to the activation energy for self-diffusion U_d only if the energy of jog formation $U_j \ll U_d$ or $U_j \gg U_d$ (in the latter case the jogs cannot be produced by fluctuation processes, but only by the intersection of dislocations, which sharply reduces the rate of climb). However, if U_j and U_d are commensurable, then $U_c > U_d$ [20].

5. The general case of dislocation creep. In the general case the processes discussed above take place in a complex dislocation structure, the elements of which may change as a result of climb or slip.

Numerous papers discuss the scheme of dislocation pile-ups suggested by Weertman [36]. If a stress σ compresses n dislocations against an obstacle, the first dislocation will be acted upon by a stress nσ. The rate of climb of this first dislocation is given by Eq. (4. 2) with the external stress σ replaced by the internal stress of the order of nσ. If the stress field produced by the obstacle is inversely proportional to the distance, then to get past the obstacle the dislocation must climb a distance inversely proportional to nσ. Hence the creep rate is proportional to $(n\sigma)^2D/kT$.

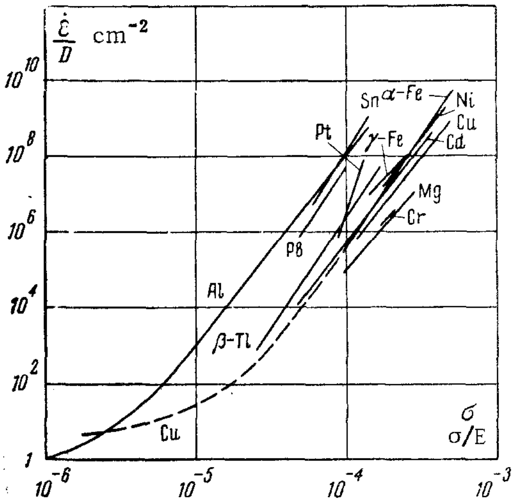

Fig. 7

If the dislocations in the pile-up come from a region of the slip plane of length L, then $n \approx 2\sigma L/gb$, and the creep rate is proportional to σ^4D/kT. Taking into account various geometrical factors leads to an additional weak dependence of $\dot{\varepsilon}$ on σ, and the final creep rate is

$$\dot{\varepsilon} = \frac{\text{const}}{kT} \sigma^m \exp \frac{-U_d}{kT} \; (3 < m < 4.5). \tag{5.1}$$

At large nσ values allowance should be made for the reduction in activation energy due to the stress energy, so that in (5. 1)

$$\sigma^m/kT \quad \text{must be replaced by} \quad \sigma^{m-2} \operatorname{sh} An\sigma\Omega/kT,$$

where A is a constant determined by the geometry of the problem.

Analysis of the results of experiments on 15 pure metals, carried out at temperatures exceeding 0. 5 T_* [37], shows that in $\ln(\dot{\varepsilon}/D)$, $\ln(\sigma/E)$ coordinates (Fig. 7) parallel straight lines with a slope $m \approx 5$ are obtained. The author believes this to be a confirmation of Eq. (5. 1). Similar results were obtained by McLean [38] in tests on creep in α iron at 550 and 700°C (m = 5. 4); in fact, he found that the dislocation density determined by direct electron microscopy increased only slightly with stress, so that the strong dependence of the creep rate on stress must have been due to the effect of the stresses on the average rate of motion of the dislocations.

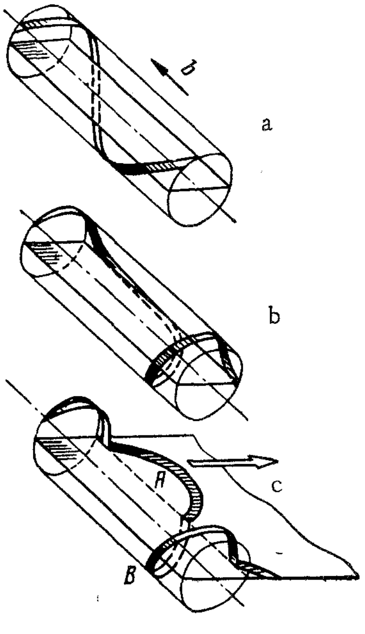

Fig. 8.

A similar investigation of the kinetics of high-temperature creep in a number of fcc metals (including single crystals) was carried out by Pines et al. [5, 9, 39]. At near-melting temperatures the relation between creep rate and stress is linear, as long as σ does not exceed a certain critical value.

When $\sigma > \sigma_0$

$$\dot{\varepsilon} = \dot{\varepsilon}_1 + \dot{\varepsilon}_2 \qquad (\dot{\varepsilon}_1 = A_1\sigma, \; \dot{\varepsilon}_2 = A_2(\sigma - \sigma_0)^m)$$

With increasing temperature, A_1 and m also increase, while σ_0 decreases. The constant A_1 depends strongly on the structure. When $\sigma < \sigma_0$ deformation is not accompanied by visible traces of slip. When $\sigma > \sigma_0$ the surface of the deformed single crystal displays a relatively small number of coarse slip bands. For aluminum the variation in σ_0 with temperature was proportional to the shear modulus G, and the creep rates obtained at various temperatures and stresses fitted a universal curve of the type:

$$\dot{\varepsilon} = Df(\sigma / G). \tag{5.2}$$

In accordance with Sherby's results [37]. For nickel the ratio σ_0/G depended on T and the creep rates failed to satisfy Eq. (5.2).

In [39] Pines and Sirenko examined the relation between σ and T and the variation in $\Delta\sigma/\Delta T$ for the case of a steady-state creep. For copper at 600-1000°C the experimental data agreed with Eq. (5.1) for m varying from 1 to 4. Pines also assumes that even in the range $T < 0.5\, T_*$ [4] taking into account the dependence of the coefficient of the exponent in (4.1) on T and σ can lead to agreement between the activation energy for creep and the energy of self-diffusion.

More recently, the mechanism of high-temperature creep in metals has been investigated by Utevskii et al. [40, 41], using the methods of transmission electron microscopy. They found that the dislocation structure of nickel and its alloys is characterized by a large number of three-dimensionally curved dislocations incapable of pure slip. The main forms of dislocations include helicoids due to the combining of point defects and screw dislocations (Fig. 8a) and further distorted by the applied stresses (Fig. 8, b, c).

Motion of the helicoidal dislocation as a whole means that the extra planes must grow longer on one side of the turn and shorter on the other. This process can be effected by diffusion flow within the turns of the helicoid. The rate of motion of a dislocation under a stress σ is calculated by analogy with (4.3) [40]:

$$v = \frac{4\pi D}{b\delta}\,\frac{\sigma\Omega}{kT}, \qquad \delta = \left(\frac{R}{H}\right)^2 \left(\ln\frac{R}{r_0} - \frac{H}{\sqrt{R^2+H^2}}\ln\frac{R^2+H^2}{Rr_0}\right). \tag{5.3}$$

Here δ is determined by the geometry of the helicoid (the radius of the turns R and the pitch H).

Numerical evaluations give a quantitative agreement between the theoretical rate of motion of the helicoidal dislocations and the value estimated from (3.1) on the basis of experimental data on the creep rate.

Judging from the results of [41], high-temperature creep is mainly determined by the non-conservative motion of dislocations due to the diffusion displacement of point defects over relatively short distances.

6. Importance of surface defects (block and grain boundaries). The joint movements of dislocations that form the boundaries of blocks also result in plastic deformation. For a crystal consisting of blocks of size L, misoriented on the average by an angle θ, the rate of deformation is

$$\dot{\varepsilon} = \theta v^{\circ}/L \tag{6.1}$$

(v° is the rate of motion of the boundaries).

The integral displacement of small-angle boundaries at increased temperatures has been observed in zinc [42]. Because of the structure of the dislocation networks in crystals, their integral motion due to an external stress is such that some of their elements must climb. The rate of displacement of a mechanically stable network in an fcc lattice controlled by the climb has been calculated in [43].

A more important factor in the mechanism of creep is the passive role played by the inter-block boundaries. In the case of diffusion creep, this reduces to the absorption and release of vacancies (with simultaneous climb of the dislocations) [4], while in the case of dislocation creep the same boundaries have a hardening effect.

The primary stage of creep usually establishes a certain equilibrium block structure. The average misorientation and the dimensions of the blocks depend on the initial structure of the material (degree of work-hardening, grain size). Other things being equal, the blocks will be larger, the higher the temperature and lower the stresses. According to several authors the block structure does not undergo any marked changes during secondary creep [28, 29]. Direct electron-microscope investigations of the dislocation structure of various metals during creep are described in papers [29, 38, 41]. In nickel at 500°C [41] and in aluminum at room temperature [29] subboundaries are formed. These are more regular than those obtained after cold working, but less perfect than the boundaries obtained after polyganization. For example, Fig. 9 [29] shows the walls of dislocations associated with creep in aluminum (experimental temperature 20°C, magnification 53 000). The structure of α-iron also shows subboundaries formed during creep [38]. The role played by grain boundaries in diffusion creep was considered in §2. When creep is caused by the motion of dislocations in the interior of the grains, the boundaries represent an obstacle. In order to ensure compatibility of deformation in adjacent grains it is necessary for slip to take place in several systems, at least in the boundary region of one of the grains.

At low temperatures, multiple slip is made difficult; therefore the boundaries harden the material, and a decrease in grain size reduces the creep rate. At high temperatures thermal activation facilitates the mutual intersection of dislocations in multiple slip. A lowering of the activation barriers for diffusion, indicated by an increase in the rate of diffusion in the boundary zone, also causes the dislocations to climb. As a result, at elevated temperatures multiple slip is facilitated, and the creep rate rises. This is in agreement with the results of an experimental investigation into the effect of the orientation of aluminum single crystals on creep [44, 45]. The simultaneous deformation of the near-boundary regions in several slip systems increases the creep rate, the more sharply the finer the grain. Nonuniform plastic deformation in the near-boundary regions of adjacent grains results in the formation of macroscopic jogs, which have a nonuniform rate of growth along the boundary and may vary considerably with time [46].

Fig. 9.

7. Creep and fracture. The creep process is usually accompanied by the initiation and growth of cracks and micropores (e. g., [47]), which raises the question of the effect of local failure on the creep rate. According to Zhurkov [27], the effect of the fracture rate on the creep rate follows from the coincidence of the parameters describing the dependence of life and creep rate on stress and temperature. Orlov [48] assumes that the development of microcracks can facilitate the passage of plastic deformation owing to the flow of dislocations into the crack.

According to other authors, the creep rate determines the fracture rate rather than vice versa. The creep process merely prepares the material for failure as a result of the accumulation of certain irreversible changes. For example, the vacancies produced in plastic deformation coagulate into pores along the grain boundaries and the growth of these pores eventually causes fracture [49]. The subordinate part played by fracture in secondary creep is also indicated by the fact that the creep rate in the tertiary stage is determined by the normal stresses and in the secondary stage by the shear stresses [50].

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