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THE KINETICS OF DISLOCATION CLIMB OVER HARD PARTICLES—II. EFFECTS OF AN ATTRACTIVE PARTICLE–DISLOCATION INTERACTION

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Abstract—A model for the high-temperature creep strength of dispersion strengthened materials is presented. It is based on part I of this paper which treated dislocation climb over "non-interacting" particles. The present part II deals with the kinetics of dislocation climb over hard particles which exert an attractive force on the dislocations. The attraction is shown to affect the climb and bypass process in two ways: (i) above a certain stress, local climb becomes a stable mechanism, and (ii) a strong threshold stress appears for detachment of the dislocation from the particle. Thermal activation of the detachment process and particle dragging are incorporated and shown to affect the threshold stress for very small particles. Qualitative comparison of predictions with the creep behaviour of dispersion-strengthened materials suggests that an attractive interaction is probably the key to a better understanding of the highly stress-sensitive creep rate in these materials.

Résumé—Nous présentons un modèle de la résistance au fluage à haute température de matériaux durcis par dispersion. Les bases de ce modèle ont été exposées dans la première partie de l'article qui traitait de la montée des dislocations en présence de particules "non interactives". Cette seconde partie concerne la cinétique de la montée des dislocations en présence de particules dures qui exercent une force d'attraction sur les dislocations. Nous montrons que l'attraction affecte le mécanisme de montée et de franchissement de deux façons: (i) au-dessus d'une certaine contrainte, la montée locale devient un mécanisme stable, et (ii) une contrainte-seuil élevée apparaît pour détacher la dislocation de la particule. Nous incorporons l'activation thermique du mécanisme de détachement et le traînage de la particule, et nous montrons que tous deux agissent sur la contrainte-seuil dans le cas de particules très petites. Une comparaison qualitative des prédictions et du comportement en fluage de matériaux durcis par dispersion suggère qu'une interaction attractive est sans doute la clef d'une meilleure compréhension de la forte influence de la contrainte sur la vitesse de fluage dans ces matériaux.

Zusammenfassung—Ein Modell für die Hochtemperatur-Kriechfestigkeit dispersionsgehärteter Legierungen wird vorgelegt. Es beruht auf dem Versetzungsklettern über "nicht-wechselwirkende" Teilchen, welches im Teil I dieser Arbeit behandelt worden ist. Der vorliegende Teil II behandelt die Kinetik des Versetzungskletterns über harte Teilchen, die eine attraktive Kraft auf die Versetzung ausüben. Diese attraktive Wechselwirkung beinflußt den Kletter- und Passierprozeß auf zweierlei Weise: (i) oberhalb einer bestimmten Spannung wird lokales Klettern ein stabiler Prozeß, und (ii), es erscheint eine kräftige Schwellspannung für das Ablösen der Versetzung vom Teilchen. Die thermische Aktivierung dieses Ablöseprozesses und der Reibungseffekt durch das Teilchen werden mit berücksichtigt; sie beeinflussen die Schwellspannung bei kleinen Teilchen. Aus dem qualitativen Vergleich der Voraussagen mit dem Kriechverhalten dispersionsgehärteter Werkstoffe geht hervor, daß eine attraktive Wechselwirkung wahrscheinlich der Schlüssel zu einem besseren Verständnis der höchst spannungsempfindlichen Kriechrate dieser Werkstoffe ist.

1. INTRODUCTION

Incorporation of incoherent, non-shearable dispersoids is an efficient way of improving the mechanical properties of metallic materials at very high temperatures. Such dispersion strengthened alloys show a drastic reduction in creep rate with diminishing stress, see e.g. [1–7]. Because stress exponents for the creep rate lie in the range $n = 20$–40, the creep behaviour of such alloys is best described by specifying a so-called "threshold stress" below which the rate of creep deformation is negligible. Threshold stresses lie typically between 0.4 and 0.8 of the Orowan stress, as measured by relatively fast deformation or estimated using standard equations.

In part I of this paper [8] it was shown that such a behaviour cannot be modelled in terms of dislocation climb over particles under equilibrium conditions. The reason is that local climb, which has been the basis of earlier models, is never stable under these conditions: a dislocation climbing over a "non-interacting" particle unravels its sharp curvature at the point where it meets the particle, until climb at the particle can proceed. This lowers the necessary line length increment, and leads to an unrealistically small threshold stress.

An explanation for the creep behaviour of dispersion strengthened materials must therefore be sought in an additional mechanism. One possibility is the existence of an attractive interaction between the climbing dislocation and the particle. Such an effect would be expected only for elastically soft particles or voids; but recent theoretical studies, see e.g. [9], have shown that at high temperatures an incoherent particle can behave elastically just like a void with an internal pressure. The reason is that diffusion can rapidly relax shear stresses and, to some extent, hydrostatic stresses imposed on the particle by the approaching dislocation. If in addition the particle–matrix interface is considered as "slipping", the dislocation can lower its energy in the vicinity of the interface. This implies that at high temperature a dislocation can be attracted towards a particle, even if it is infinitely hard [9].

Experimental evidence for an attractive interaction comes from TEM studies of dislocation configurations in crept oxide-dispersion strengthened alloys [10–12]. The dislocations are predominantly captured in a situation where they seem to stick to the "departure" side of the particles, indicating that they are immobilized by an attractive force after climb over the particle has been completed. Under weak-beam conditions, the dislocation contrast is usually well visible in (or near) the particle–matrix interface [11]. This observation suggests that the dislocation core has not relaxed completely as might be expected in the presence of a phase boundary.

That full relaxation of the dislocation is not required for the attraction effect to become the strength-determining mechanism has been established by Arzt and Wilkinson [13]. In their model, the magnitude of the threshold stress is calculated as a function of the interaction strength and compared with the threshold for local climb. The result shows that only a small interaction, corresponding to a relaxation of the dislocation line energy by about 6%, is necessary for detachment of the dislocation from the particle to become the overriding threshold mechanism.

Although the idea of an attractive interaction is thus consistent with experimental and first theoretical results, the consequences of such a mechanism are not totally clear. As long as a kinetic model for the climb process with an attractive interaction is lacking, the influence of the attraction on the geometry of climb remains unknown. The aim of the second part of this paper is therefore to establish the kinetics of dislocation climb over an attractive particle, with similar simplifying assumptions as in part I. We consider the effect of the interaction on the equilibrium profile of the dislocation at, and in the vicinity of, the particle. The resulting dislocation velocities are calculated, and the threshold stress for detachment is incorporated. In this way several features which are reflected in the creep behaviour of dispersion strengthened alloys emerge naturally from the analysis. The symbols used in the development are listed in Table 1.

Table 1. List of symbols used in parts I and II

τ	shear stress
τ_d	shear stress for detachment
τ_{th}	threshold shear stress
τ_{th}^r	threshold shear stress for restricted climb
τ_O	Orowan stress (in shear)
f	particle volume fraction
2λ	mean planar particle spacing
h, d, β	height, width and ramp angle of the particles (cf. Fig. 2 in part I)
G	shear modulus
b	magnitude of a Burgers vector
x, y, z	coordinates (cf. Fig. 2 in part I)
x_0, z_0	coordinates describing the dislocation profile (cf. Fig. 2 in part I)
x_0^*, z_0^*	dislocation profile at the transition to local climb
μ	chemical potential for vacancies
k_B	Boltzmann's constant
T	absolute temperature
A_{ABD}	area under the climbing dislocation segment
D_v	lattice diffusivity
a_pD_p	pipe cross section times pipe diffusivity ($a_p \approx b^2$)
a_v	area of a vacancy
T_{AC}, T_{CD}	line energy of the dislocation segment AC and CD resp.
ρ	radius of curvature of the climbing dislocation (cf. Fig. 2 in part I)
ρ_r	ρ for restricted climb (Fig. 5)
t_c	time for a dislocation to climb over a particle
t_c^*	$=d^3/60C_iGb^4$, normalizing parameter for t_c
n_c	fraction of particles to be climbed over
γ_0	equilibrium contact angle of the dislocation at the particle
k	relaxation factor
M	mobility of a dispersoid particle
ρ_m	mobile dislocation density

2. STABILITY OF LOCAL CLIMB DUE TO PARTICLE–DISLOCATION ATTRACTION

One effect of an attractive interaction is the modification of the equilibrium dislocation profile at, and in the vicinity of, the particle during climb. The dislocation will attempt to maximize its line length at the particle–matrix interface, even at the expense of more total line length. This aspect is treated in this section; we will turn to the effect on dislocation detachment in Section 3.

An attractive interaction between dislocations and particles can be modelled, as has been done before [13], by assigning a lower line energy T_{AC} to the dislocation segment residing in or near the particle–matrix interface

$$T_{AC} = k \cdot T_{CD} \tag{1}$$

T_{CD} is the line energy of a lattice dislocation remote from the particle; it will be approximated by

$$T_{CD} \cong \tfrac{1}{2} Gb^2 \tag{2}$$

where G is the shear modulus of the matrix and b the magnitude of the Burgers vector of a dislocation. The factor k in equation (1) can be thought of as a "relaxation factor" [13]: it describes the extent to which the dislocation relaxes its energy by interaction with the particle–matrix interface. For $k = 1$, no

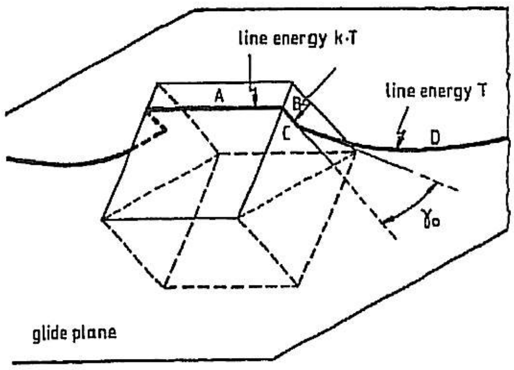

Fig. 1. Equilibrium profile of a dislocation climbing over an attractive, cube-shaped particle. The attractive interaction is modelled by assigning a reduced line tension to the dislocation segment along the interface. This leads to a stable dislocation bend at C, with equilibrium angle γ_0. Depending on k, fully or partially local climb can occur (see text).

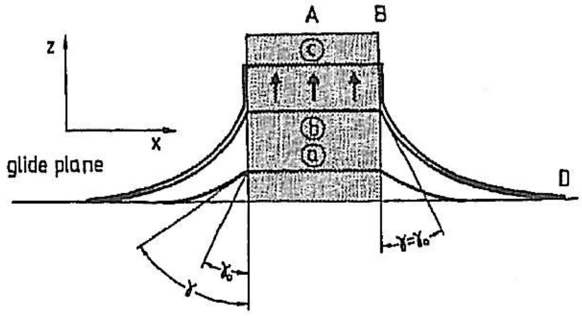

Fig. 2. Sequence of events during climb over an attractive particle: (a) initial general climb ($\gamma > \gamma_0$), (b) transition to local climb ($\gamma = \gamma_0$), and (c) local climb with constant γ.

relaxation, and thus no attractive interaction, occurs; such "non-interacting" particles have been considered in part I. $k = 0$ signifies the other extreme, i.e. complete relaxation and maximum attractive interaction, with the particle behaving as if it were a void. Relaxation, which occurs by diffusion, is a time-dependent process: its extent will therefore be determined by the ratio of its time constant to the time for climb bypass. As a result, k values which lie intermediate between 0 and 1 are possible.

Because of the discontinuity in line energy, a sharp bend of the dislocation at the particle surface can now be stable (at C in Fig. 1). The situation is analogous to two impinging surfaces of different specific energy: they form a well-defined "wetting angle" which depends on the ratio between their energies. A similar analysis gives the equilibrium contact angle γ_0 for the dislocation

$$\cos \gamma_0 = k. \tag{3}$$

For "non-interacting" particles with $k = 1$, this condition is equivalent to the requirement that a smooth tangent exists ($\gamma_0 = 0$). In this case a sharp dislocation bend is unstable, and the kinetics of climb reduces to that described in part I. For maximum interaction ($k = 0$), the dislocation will enter the particle–matrix interface perpendicular to the particle surface ($\gamma_0 = \pi/2$).

A stable sharp bend in the dislocation makes local climb a viable mechanisms. Consider the sequence of events as a dislocation climbs up an attractive particle (Fig. 2):

(a) At first, the dislocation climbs in "general" mode and its shape (x_0, z_0) is given by equation (10) of part I. The angle γ is large compared with γ_0.

(b) As the dislocation climbs further, it unravels to larger values of x_0; this leads to a decrease of the angle γ. Eventually, when this angle reaches the equilibrium value the dislocation segment CD can remain in its equilibrium position without accepting further vacancies.

(c) From now vacancies are only accepted along AB and γ_0 is retained. The rest of the particle is surmounted by local climb.

The point z_0^* at which the transition to local climb is possible can be calculated by first expressing the angle γ formed by the segments AC and CD during general climb, as a function of z_0/x_0, ramp angle β and normalized shear stress τ/τ_0 [where τ_0 is the Orowan stress given in equation (2) of part I]. Applying trigonometry we obtain after considerable rearrangement

$$\cos \gamma = \frac{\cos \beta}{1 + \left(\frac{z_0}{x_0}\right)^2} \cdot \left\{ 2\frac{z_0}{x_0} + \left(1 - \frac{z_0^2}{x_0^2}\right)\frac{\tau}{\tau_0} \tan \beta \right\}. \tag{4}$$

This angle becomes equal to the equilibrium value [equation (3)] when

$$\frac{z_0^*}{x_0^*} = \frac{1 - \sqrt{1 - \left(\frac{k}{\cos \beta}\right)^2 + (\tau/\tau_0)^2 \tan^2 \beta}}{\frac{k}{\cos \beta} + \frac{\tau}{\tau_0} \tan \beta}. \tag{5}$$

The unravelling distance x_0^* at the point of transition is given by equation 10 of part I

$$x_0^* = d \left\{ \sqrt{1 - (\tau/\tau_0)^2} \left(\frac{x_0^*}{z_0^*} + \frac{z_0^*}{x_0^*} \right) \times \left[\frac{\tau}{\tau_0} \frac{1}{\tan \beta} - 1 + \left(1 - \frac{z_0^*}{x_0^*}\right)^{9/4} \right] \right\}^{-1}. \tag{6}$$

By combining equations (5) and (6), an explicit (but lengthy) expression for z_0^* at the point of transition can be obtained. The results are plotted in Fig. 3 for different values of k and will be discussed in Section 5.

3. DETACHMENT THRESHOLD DUE TO PARTICLE–DISLOCATION ATTRACTION

Besides stabilizing local climb, an attractive particle–dislocation interaction has a second effect on bypass by a climbing dislocation: it introduces a threshold stress for detachment of the dislocation from the particle after climb has been completed. The origin of this threshold is the energy which must be supplied to substitute new dislocation line

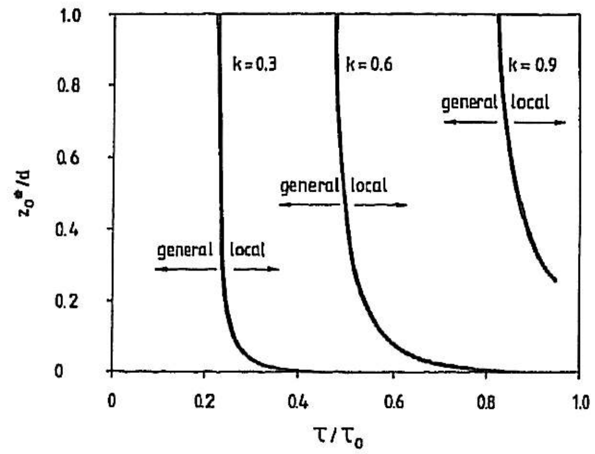

Fig. 3. Regimes of local and general climb as a function of climb height z_0 and applied stress τ/τ_0, for various relaxation parameters k.

of full energy T_{CD} for dislocation line relaxed to energy $k \cdot T_{CD}$ (Fig. 4). The magnitude of the resulting threshold for detachment is calculated in the Appendix

$$\frac{\tau_d}{\tau_0} = \sqrt{1-k^2}. \quad (7)$$

Because the critical configuration which determines the detachment threshold is reached only at the very point of detachment, the threshold is independent of particle size and shape. Therefore equation (7) is identical with the result obtained from a limit consideration in [13] for spherical particles. Also the detachment threshold applies irrespective of whether climb is local or general. It is further not affected by the position of the glide plane with respect to the particle it intersects.

The detachment threshold truncates the overall climb velocity, overriding all other considerations concerning climb kinetics. At stresses below τ_d, a dislocation remains stuck to the particle, although climb up the "arrival" side of the particle may have been fast and easy. In fact, the shape of the force–distance profiles for a dislocation climbing over a spherical particle is such [13] that the dislocation hardly "feels" the detachment stress before it actually reaches the point of detachment. (It is much like a ball rolling down a ramp with a hole at the end: the hole does not affect the speed of the ball until it finally captures it.) When calculating climb velocities as described in the following section, the detachment threshold is taken into account by truncating the velocities at the stress given by equation (7).

The detachment threshold has up to now been considered as a reliable, athermal limit to dislocation motion. There are however two ways in which this threshold could be circumvented: thermally activated dislocation detachment from the particle or particle dragging by the captured dislocations. Because both processes could potentially lower the effective threshold stress, their relative importance will be estimated in the following.

It is shown in the Appendix that thermally activated detachment of a dislocation from an attractive particle gives the appearance of a reduced threshold stress of the following form

$$\frac{\tau_d(T)}{\tau_0} = \sqrt{1-\left(k+\frac{2k_B T}{Gb^2 d}\ln\frac{\dot{\varepsilon}_0}{\dot{\varepsilon}}\right)^2} \quad (8)$$

where $\dot{\varepsilon}$ is the strain rate, and $\dot{\varepsilon}_0$ the reference creep rate (e.g. that of the particle-free material at the same stress). This expression reduces to equation (7) for $T/d \to 0$. $\tau_d(T)/\tau_d$ is plotted as a function of $(2k_B T/Gb^2 d)$ in Fig. 5, with $\dot{\varepsilon}_0/\dot{\varepsilon} = 10^{10}$ (see Appendix). The threshold is found to disappear at very high temperatures and/or small particle sizes, i.e. for

$$\frac{T}{d} \geqslant \frac{Gb^2}{2k_B \ln(\dot{\varepsilon}_0/\dot{\varepsilon})}(1-k). \quad (9)$$

A rough order-of-magnitude estimate gives as a limiting particle size $d \approx 10$ nm ($k = 0.9$, $G \approx 10^{11}$ Nm^{-2}, $b \approx 10^{-10}\,m$, $k_B T \ln(\dot{\varepsilon}_0/\dot{\varepsilon}) \approx 2 \times 10^{-19}$ Nm). For stronger interactions ($k < 0.9$), this limit is shifted to even smaller particle sizes.

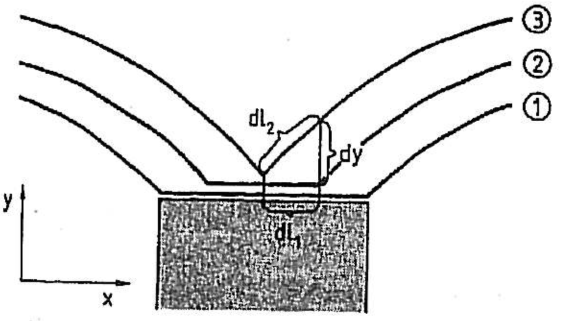

Fig. 4. Illustration of the energy balance for a dislocation at the point of detachment from the particle. The line energy is T_{CD} in the lattice, and is reduced to $k \cdot T_{CD}$ at the particle–matrix interface. A detachment threshold τ_d arises because dislocation line length of full line energy has to be substituted for relaxed dislocation line.

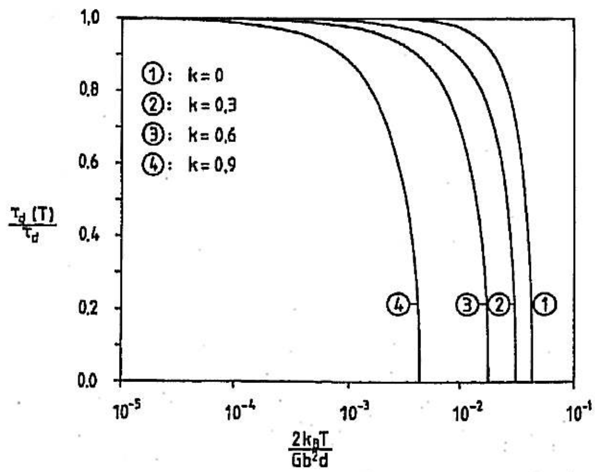

Fig. 5. Normalized shear stress vs the parameter $2k_B T/Gb^2 d$ for $\dot{\varepsilon}_0/\dot{\varepsilon} = 10^{10}$, according to equation (8). Thermal activation contributes significantly to the detachment process from small particles with diameter $d < (0.02 k_B T)/(Gb^2)$.

Another possibility for a captured dislocation to continue moving is by pulling the particle in its wake. The resulting strain rate is derived in the Appendix

$$\dot{\varepsilon}_{\rm pm} = \frac{2\rho_{\rm m} M b T_{\rm CD}[\tau/\tau_0]}{1 + \dfrac{\varepsilon}{2\rho_{\rm m} b \lambda}} \tag{10}$$

where $\rho_{\rm m}$ is the mobile dislocation density, M the mobility of the particle, and ε the accumulated strain.

Deformation with particle movement is a process without a threshold stress and is therefore expected to dominate at low stresses. However, it is self-exhausting: its rate decreases with strain because the dislocation collects particles as it sweeps through the material. Therefore particle dragging cannot contribute to steady-state deformation, it can only give rise to a transient creep rate. The maximum strain due to this mechanism is usually negligible: it is estimated in the Appendix to be typically 10^{-4}. Indirect evidence for particle dragging is found in TEM micrographs showing agglomerates of fine particles in a dispersion strengthened superalloy after creep [19].

4. THE KINETICS OF COMBINED GENERAL + LOCAL CLIMB WITH A DETACHMENT THRESHOLD

The forward velocity of a dislocation which climbs over a particle has been expressed in part I as

$$\frac{{\rm d}y}{{\rm d}t}(z_0) = C_l \frac{|\mu|}{|{\rm d}A/{\rm d}y|} \tag{11}$$

where A is the projected area under the climbing dislocation segment, and C_l is a kinetic constant which depends on whether volume diffusion or lattice diffusion is faster in supplying vacancies to the climbing dislocation segment.

The chemical potential μ for vacancies along the climbing dislocation and the area increment ${\rm d}A/{\rm d}y$ depend on whether or not the transition to local climb has already occurred. In the general climb mode, μ is unaffected by the attractive interaction, because the segment at the interface has constant length. The chemical potential is then given as before by equation (6) of part I. Once the transition has taken place, the chemical potential for addition of further vacancies along AB changes. In the Appendix it is calculated, with the following result

$$\mu_{\rm AB}^{\rm loc} = \frac{2a_{\rm v}}{{\rm d}\sin\beta\cos\beta}(T_{\rm AC}\sin\beta - \tau\lambda b). \tag{12}$$

Unlike the potential for general climb, that for local climb is independent of z_0.† The increment in the area under the climbing dislocation during an infinitesimal advance is given by a simple expression

$$\frac{{\rm d}A}{{\rm d}y} = \frac{{\rm d}}{2}\sin\beta\cos\beta \tag{13}$$

The dislocation velocity during local climb can now be obtained explicitly by combining equation, (11), (12) and (13)

$$\frac{{\rm d}y}{{\rm d}t_{\rm loc}} = C_l \frac{4a_{\rm v}}{d^2\sin^2\beta\cos^2\beta}(\tau\lambda b - T_{\rm AC}\sin\beta). \tag{14}$$

We arrive at the total time for a dislocation to bypass a particle by combined general + local climb after appropriate integration

$$t_{\rm c} = \int_0^{y^*} \frac{1}{({\rm d}y/{\rm d}t)}{\rm d}y + \int_{y^*}^{y_{\rm max}} \frac{1}{({\rm d}y/{\rm d}t)_{\rm loc}}{\rm d}y \tag{15a}$$

where y^* denotes the y-coordinate at the point of transition z_0^* given by equations (5) and (6). To incorporate the detachment threshold in a simple way, we require further

$$t_{\rm c} = \infty \quad \text{for} \quad \tau < \tau_{\rm d}. \tag{15b}$$

Using equations (15), dislocation velocities (i.e. reciprocals of the bypass times) were calculated as a function of applied stress for several particle parameters and values of k. Some results are shown in Figs 6 and 7.

5. DISCUSSION

The new element which this paper adds to the discussion on dislocation climb over particles is the

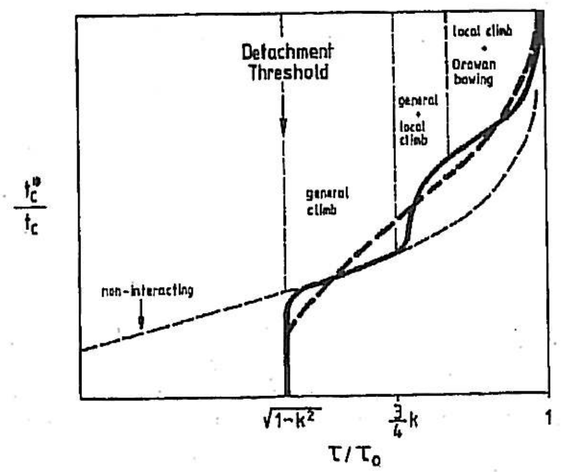

Fig. 6. Dislocation velocity t_c^*/t_c as a function of normalized shear stress, τ/τ_0 (schematically). The regions of the different climb variants are indicated approximately. For comparison, the result for "non-interacting" particles ($k = 1$) is also included (fine dashed line—see part I [8]).

†At the point of transition given by equation (5) also the driving force for local climb [equation (12)] starts to exceed that for general climb [equation (6) in part I]. Therefore the transition could have been based alternatively on a driving force criterion. Because of the simplifications in dislocation geometry introduced in part I, the driving force criterion diverges somewhat from the equilibrium angle criterion for $k > 0.8$. We have chosen to use always the angle criterion for describing the point of transition.

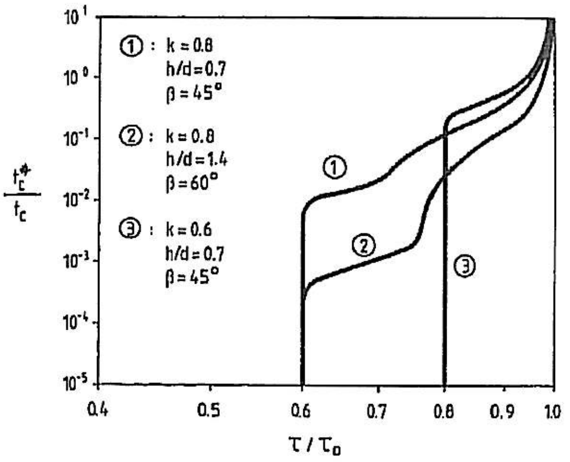

Fig. 7. Dislocation velocities vs normalized shear stress for different relaxation parameters k and particle geometries.

assumption of an attractive interaction between particles and dislocations. The present model describes, to a certain approximation, the details of the bypass process of a dislocation past an attractive particle. It hinges on the assumption that the dislocation can adopt a minimum energy configuration; this may not always be the case, but this assumption appears to be better justified for a slow high-temperature process than the *a priori* postulate of local climb. It has been shown that an interaction has two separate effects on the climb and bypass process:

(i) it can stabilize local climb during certain parts of the climb process, and

(ii) it introduces a strong threshold stress for dislocation detachment from the particle.

The reason for the transition from general to local climb lies in the assumption that the dislocation always adopts a configuration of minimum energy. For the simple particle shape assumed in this paper, several observations can be made. For a given interaction strength, the tendency towards local climb depends on the applied stress. In Fig. 3 the regimes of the two climb variants "general" and "local" are demarcated as a function of stress. At sufficiently high stresses, climb is local from the beginning; the borderline is obtained by requiring $z_0^* = 0$ in equation (5) (which determines where the transition occurs), i.e.

$$\frac{\tau}{\tau_0} > \frac{k}{\sin\beta} \tag{16}$$

In this range, the previous models for local climb would in principle be applicable. However, "non-interacting" particles ($k = 1$)—the only ones considered in previous models—would require a stress at least as high as the Orowan stress in order to ensure local climb. When attractive particles are considered ($k < 1$), then the borderline to purely local climb falls below the Orowan stress; a stress "window" then opens up in which the climb part of the present model reduces to the former models for local climb. But at the same time the detachment threshold appears, and it is usually the dominating threshold.

At sufficiently low stresses, the equilibrium contact angle is never reached and climb is general throughout. This situation occurs when the square root in equation (5) becomes unreal, i.e. when

$$\frac{\tau}{\tau_0} < \sqrt{\left|\left(\frac{k}{\sin\beta}\right)^2 - \frac{1}{\tan^2\beta}\right|}. \tag{17}$$

For $k < 0.8$, a simple, approximate borderline for equiaxed particles of height d and ramp angle 45° is obtained by inspection of Fig. 3

$$\frac{\tau}{\tau_0} \lesssim \frac{3}{4}k. \tag{18}$$

Finally, in the intermediate range a transition from initially general to local climb occurs.

The detachment threshold is the second, and usually more important, effect of an attractive interaction. The treatment of thermally activated detachment, leading to equation (8), shows that it is not a true athermal threshold; it disappears when the length over which the dislocation adheres to the dispersoid reaches the magnitude of a few nm. In dispersion strengthened alloys the average dispersoid size is commonly larger than this, but the smaller particles in the distribution will cease to be effective obstacles. Hence this mechanism defines a potential limit to the fineness of an efficient particle dispersion—it is an "intrinsic" limitation to the dispersion strengthening achievable by a given volume fraction of dispersoid. When the particles are spherical the "detachment length" can be much smaller than the particle diameter. This provides a possible explanation for the temperature dependence of the threshold stress, as will be detailed elsewhere [20].

The theoretical dislocation velocities reflect the existence of the different climb and bypass mechanisms (Fig. 6). At low velocities, the detachment threshold is the dominating feature. Just above this threshold, general climb occurs, with a dislocation velocity identical to that for climb over a "non-interacting" particle (fine dashed line in Fig. 6); the kinetics of general climb is not affected by the interaction. At $\tau/\tau_0 \approx 3/4\,k$, where the transition to local climb occurs, the stress sensitivity increases: the reason is that the point of transition is highly stress-dependent (Fig. 3) and local climb, which requires fewer vacancies, is kinetically faster than general climb. The magnitude of this "bump" depends on particle geometry and increases with the aspect ratio h/d and the ramp angle β. Finally, at still higher stresses, a transition to Orowan bowing occurs and the behaviour of the material eventually approaches that of the dispersoid-free material.

Figure 7 illustrates the effects of different values of k and particle parameters h/d and β. The particle geometry influences the climb velocity above the threshold substantially. At $k = 0.6$ the detachment threshold exceeds the stress at which the transition to local climb occurs, and the general climb region disappears altogether.

The main weakness of the model as it stands lies in the fact that only a single particle of simple shape is considered. Therefore an abrupt transition to local climb is predicted, resulting in the peculiar features shown in Fig. 6, which are not reflected in experimental creep data. The problem is that the fine details of the climb kinetics are very sensitive to particle shape and will remain difficult to model for realistic particle geometries. It can be expected that the oscillations in Fig. 6 will disappear when an ensemble of statistically distributed obstacles with varying h/d and β is considered, resulting possibly in the heavy dashed line in Fig. 6. A computer simulation will be necessary to verify this point.

Overall, the assumption of an attractive interaction seems to have great potential for explaining the creep behaviour of dispersion strengthened materials. The high stress sensitivity at low creep rates can be attributed, at least qualitatively, to a detachment threshold and—depending on particle geometry, but usually less important—to the transition to local climb. The relaxation parameter k plays therefore a central role for the creep strength; it is not clear whether and how its value depends on the microstructure, e.g. dispersoid size, structure of the dispersoid–matrix interface etc. To allow a quantitative prediction of threshold stresses, a more detailed understanding of the microscopic interaction mechanisms between dislocations and incoherent particles at high temperature is therefore urgently needed.

6. SUMMARY

1. The high stress sensitivity of the creep rate in dispersion strengthened alloys can be explained in terms of a threshold stress for detachment of dislocations captured at attractive dispersoid particles. The assumption of such an interaction is indispensable because "non-interacting" particles of low volume fraction cannot suppress creep (see part I).

2. Besides introducing a detachment threshold, an attractive interaction also stabilizes local climb, leading to an increased creep rate at higher stresses. Compared to the detachment threshold, this is an effect of secondary importance (unless the particles have an unusually high aspect ratio and/or ramp angle).

3. Thermally activated escape of captured dislocations and particle dragging can lower the threshold stress for very small particles. This sets a theoretical limit to the fineness of an efficient particle dispersion at a particle diameter of a few nm.

4. The attractive interaction is probably due to dislocation relaxation in the vicinity of the particle–matrix interface. It arises only in the presence of incoherent particles. Clarification of the detailed atomistic process at these interfaces requires further work.

5. The mechanism of dispersion strengthening as it is proposed here is consistent with TEM observations of dislocation structures in ODS superalloys. The theoretical considerations suggest that it is probably generally applicable to other dispersion strengthened systems.

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REFERENCES

1. R. W. Lund and W. D. Nix, *Acta metall.* **24**, 469 (1976).
2. J. Lin and O. D. Sherby, *Res Mechanica* **2**, 251 (1981).
3. R. Petkovic-Luton, D. J. Srolovitz and M. J. Luton, *Frontiers of High Temperature Materials* II, p. 73. INCOMAP, New York (1983).
4. A. H. Clauer and N. Hansen, *Acta metall.* **32**, 269 (1984).
5. R. C. Benn and S. K. Kang, in *Superalloys 1984* (edited by M. Gell, C. S. Kortovich, R. H. Bricknell, W. B. Kent and J. F. Radavich), p. 321. T.M.S.-A.I.M.E., Warrendale, Pa (1984).
6. J. Peterseim and G. Sauthoff, *Steel Res.* **57**, 19 (1986).
7. R. F. Singer and E. Arzt, in *High Temperature Alloys for Gas Turbines and Other Application* (edited by W. Betz, R. Brunetaud, D. Coutsouradis, H. Fischmeister, T. B. Gibbons, I. Kvernes, Y. Lindblom, J. B. Marriott and D. B. Meadowcroft), p. 97. Reidel, Dordrecht (1986).
8. J. Rösler and E. Arzt, *Acta metall.* **36**, 1043 (1988).
9. D. J. Srolovitz, M. J. Luton, R. Petkovic-Luton, D. M. Barnett and W. D. Nix, *Acta metall.* **32**, 1079 (1984).
10. V. C. Nardone and J. K. Tien, *Scripta metall.* **17**, 467 (1983).
11. J. H. Schröder and E. Arzt, *Scripta metall.* **19**, 1129 (1985).
12. E. Arzt and J. H. Schröder, in *High Temperature Alloys for Gas Turbines and Other Application* (edited by W. Betz, R. Brunetaud, D. Coutsouradis, H. Fischmeister, T. B. Gibbons, I. Kvernes, Y. Lindblom, J. B. Marriott and D. B. Meadowcroft), p. 1037. Reidel, Dordrecht (1986).
13. E. Arzt and D. S. Wilkinson, *Acta metall.* **34**, 1893 (1986).
14. L. M. Brown and R. K. Ham, *Strengthening Methods in Crystals* (edited by A. Kelly and R. B. Nicholson), p. 9. Elsevier, Amsterdam (1971).
15. R. W. Weeks, S. R. Pati, M. F. Ashby and P. Barrand, *Acta metall.* **17**, 1403 (1969).
16. R. O. Scattergood and D. J. Bacon, *Acta metall.* **30**, 1665 (1982).
17. M. F. Ashby, *Proc. 1st Int. Symp. Metall. Mater. Sci.*, Risø edited by N. Hansen, A. R. Jones and T. Leffers, p. 1 (1980).
18. E. Arzt, in *Handbook of Metallic Composites* (edited by T. Glasgow and J. D. Whittenberger). To be published.
19. J. H. Schröder, Doctoral Dissertation, Univ. Stuttgart (1987).
20. E. Arzt and J. Rösler. To be published.

APPENDIX

Calculations

For calculating the *detachment threshold* [equation (7)], consider the energy balance of an advancing dislocation at the point of detachment from the particle (Fig. 4). We equate the work done during an infinitesimal advance by dy with the energy required for exchanging a dislocation segment of length dl_1 and energy kT_{CD} by a segment of length dl_2 and energy T_{CD}

$$\tau b\lambda \, dy = T_{CD}(dl_2 - k \, dl_1). \quad \text{(A1)}$$

Using

$$(dl_2)^2 = (dl_1)^2 + (dy)^2 \quad \text{(A2)}$$

and

$$dl_2 = dy \frac{\tau_O}{\tau} \quad \text{(A3)}$$

we get for the detachment threshold

$$\tau_d = \tau_O \sqrt{1 - k^2}. \quad \text{(A4)}$$

Thermal activation of dislocation detachment from an attractive particle is modelled as follows. Consider a dislocation captured at the departure side of a particle of width d, as shown in Fig. 4. If the applied stress is insufficient to allow detachment from the particle ($\tau < \tau_d$), then in order to escape by thermal activation the dislocation will have to surmount a potential hill of magnitude

$$U \approx \frac{Gb^3}{2}\left(\frac{d}{b}\right)[\sqrt{1-(\tau/\tau_O)^2} - k]. \quad \text{(A5)}$$

This result obtained by subtracting the left-hand side of equation (A1) from the right-hand side and multiplying by $\Delta l_1 = d$. The potential hill vanishes for $\tau = \tau_d$ [equation (7)] as required.

The resulting strain rate can be written phenomenologically as

$$\dot{\varepsilon} = \dot{\varepsilon}_0 \exp\left(-\frac{U}{k_B T}\right) \quad \text{(A6)}$$

where $\dot{\varepsilon}_0$ is the equivalent strain rate of the particle-free material. Combining equations (A5) and (A6), and solving for τ/τ_O gives equation (8) in the text. Strictly speaking it does not define a threshold stress, but the stress necessary for achieving a certain strain rate ratio $\dot{\varepsilon}_0/\dot{\varepsilon}$ (as a function of k, T and d).

Equation (A6) can however give the appearance of a threshold stress when the resolution limit for measuring strain rates is considered. As a realistic example we set $\dot{\varepsilon} \approx 10^{-12}$/s and take for $\dot{\varepsilon}_0$ a typical strain rate value for the particle-free material, at the threshold stress for the particle-containing alloy (see e.g. [1]: $\dot{\varepsilon}_0 \approx 10^{-2}$/s). For plotting Fig. 5 we therefore use $\dot{\varepsilon}_0/\dot{\varepsilon} = 10^{10}$. Because this ratio enters equation (8) only logarithmically, its choice is not very critical.

Another possibility for continued dislocation motion at $\tau < \tau_d(T)$ is the *dragging of particles* by the dislocation captured on the "departure" side. The velocity of the particle is given by

$$v = M \cdot F \quad \text{(A7)}$$

F is the force exerted on the particle by the dislocation captured on the detachment side

$$F = 2T_{CD}[\tau/\tau_O]. \quad \text{(A8)}$$

If diffusion along the particle–matrix interface (with diffusivity D_i) is considered to be rate-limiting, then the particle mobility M is given, for cubes of edge length d, by

$$M = \frac{4\delta D_i \Omega}{k_B T d^4}. \quad \text{(A9)}$$

Similar expressions apply when other processes control the rate [17].

The resulting creep rate due to particle movement is

$$\dot{\varepsilon}_{pm} = \rho_m b v \quad \text{(A10)}$$

where ρ_m is the mobile dislocation density. It is important however that deformation with particle movement is a self-exhausting process: the dislocation collects particles on its way, such that their spacing along the dislocation line becomes

$$\lambda' = \frac{\lambda}{1 + \dfrac{\varepsilon}{2\rho_m b\lambda}}. \quad \text{(A11)}$$

The resulting strain rate decreases with strain ε as is evident from equation (10) in the text. The limiting strain ε_{lim} that can be achieved by this mechanism is reached when the particle spacing along the dislocation is equal to their size ($2\lambda' = d$)

$$\varepsilon_{lim} = \left(\frac{2\lambda}{d} - 1\right)2\rho_m b\lambda. \quad \text{(A12)}$$

A typical value for small volume fractions ($2\lambda/d \approx 10$) is $\varepsilon_{lim} \approx 10^{-4}$ (setting $\rho_m = 10^{12}/m^2$, $b = 10^{-10}$ m, $\lambda = 10^{-7}$ m). Therefore the contribution of the particle dragging mechanism can probably be neglected in most cases.

The *chemical potential* μ_{AB}^{loc} for local climb [equation (12)] is obtained as follows. We assume, as in [14], that the segment along the side face of the particle (BC in Fig. 1) forms a right angle with the ramp contour. Then the line length increment during a dislocation advance by dy is simply

$$\frac{dl}{dy} = \sin\beta \quad \text{(A13)}$$

and the increment in area under AB is

$$\frac{dA_{AB}}{dy} = \frac{d}{2}\sin\beta\cos\beta. \quad \text{(A14)}$$

Substituting these expressions in the general equation for μ [equation (A1) in part I] leads to equation (12) in the text.