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The 1962 Horace W. Gillett
Memorial Lecture

on

" Progress in Understanding High-Temperature Creep "

by

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The 1962 Horace W. Gillett Memorial Lecture

PROGRESS IN UNDERSTANDING HIGH-TEMPERATURE CREEP

It is with sincere personal gratitude to a great American Metallurgist that I welcome this opportunity to honor the memory of Horace Wadsworth Gillett. His extensive contributions to the advancement of engineering constitute finely fashioned monuments to a life of service devoted to the benefit of all mankind.

Dr. Gillett was a modest man of broad interests and many attainments. It was one of these special achievements for which I remember him best. Dr. Gillett, in the informal comfort of his smoke-filled office at Battelle, was an outstanding teacher. I had the good fortune to be a Post-Doctoral Fellow at Battelle from 1936 to 1938. And during this time I, among many others, took advantage of his willingness to guide young men in their search for knowledge. The door to his office was always open. And, once the problem was posed, his fertile mind had several imaginative approaches and unique solutions to suggest. He was a gentle counselor who took time to discuss the intimate details of a problem but he was also a most severe critic whose impatience with technical incompetence was reflected in retorts from his sharp, witty tongue. Each time I left his office I felt I was a wiser man.

We honor Horace Wadsworth Gillett for the man he was; and today I wish especially to pay tribute to him as a great educator in the highest sense of that title he so well earned.

Dr. Gillett's technical interests covered the entire gamut of metallurgy. He established one of the first creep testing laboratories in these United States; and throughout his life he remained active in this field. It is, therefore, fitting that this lecture be devoted to the subject of high-temperature creep of metals.

1. The Experimental Basis for Understanding Creep

The time-dependent plastic deformation of materials under constant stress is known as creep. Creep can take place only as a result of thermal activation of deformation processes. Consequently, the shear strain rate of creep, $\dot{\gamma}$, for the i th mechanism can be written formally in terms of the probability of a thermal fluctuation[1] as

$$\dot{\gamma}_i = f_i\{\tau, st, T\} e^{-\frac{g_i(\tau, st, T)}{kT}} \tag{1}$$

where k is the Boltzmann constant and T is the absolute temperature. Both the frequency factor f_i which represents the strain per second per successful fluctuation, and the free energy of activation g_i may depend on the applied shear stress, τ, the instantaneous structure of the material, st, and the temperature, T. Occasionally, several mechanisms operate simultaneously. If they operate in sequence, such that the second can take place only after the first mechanism has occurred, etc., the $\dot{\gamma}_i$ of Eqn. (1) must, under steady state conditions, as in all cases of chemical reaction rates, refer to the slowest step in the sequence. When several independent non-sequential mechanisms operate, however, the net creep rate is $\dot{\gamma} = \sum \dot{\gamma}_i$ where the summation is taken over all of the independent mechanisms. Frequently, over certain regions of stress, structure, and temperature, the creep rate arises principally from a single mechanism.

Undoubtedly, the most elusive quantity to define precisely in the phenomenological equation for the creep rate is the structure factor, st. By structure I mean all atomic and molecular arrangements including those involving the kind, number, and distribution of dislocations, point imperfections, grain size and orientations that are significant to creep. To illustrate the importance of this factor I merely refer to Fig. 1.[2] Since the stress and temperature were held constant, the variation of creep rate with time must be ascribed to the variation in structure as the test progressed. Occasionally, the changes in structure during creep will reduce the creep rate by a factor

of 10^6 or more during a single test.

The objective of fundamental investigations on creep is the determination of the rate-controlling mechanisms. When the mechanisms of creep are uncovered in sufficient detail, it should be possible to predict the shape of the creep curve and thus facilitate creep testing as well as design for creep resistance. Furthermore, when the mechanisms are known it may be possible to introduce factors that block the prominent mechanisms and thereby reduce the creep rate.

The data of creep strain as a function of time for fixed conditions of stress and temperature, as obtained in ordinary creep tests, are generally unsuited for uncovering creep mechanisms. This arises because the basic creep equation, Eqn. (1), does not contain the explicit relationship of strain as a function of time. During the course of creep the structure factor, $\mathcal{S}t$, changes. As shown in Fig. 2, however, the rate of change of the structure is more rapid with strain for tests at higher stresses.[3] Furthermore, when the stress is changed abruptly from 4000 to 2000 psi, the creep rate at A is less than that obtained upon creeping exclusively at 2000 psi to the same strain at B. Since the externally controllable variables of stress, temperature and strain are identical at A and B whereas the creep rates are different, we must conclude that the substructure depends on the details of the stress path. Consequently, the creep rate cannot be expressed in terms of a mechanical equation of state that implies that the creep rate is exclusively a function of the instantaneous values of the stress, temperature, and strain. After extensive creep under the same stress, as shown at points C and D, the same creep rate is obtained regardless of the early differences in stress history. This is typical of all high temperature creep of pure materials and results from the balance between the rate of strain hardening and the rate of recovery that takes place during the constant secondary creep rate. Thus, a constant structure, $\mathcal{S}t$,

prevails over the secondary stage of creep. But a more creep-resistant structure is developed during the secondary stage of creep under a higher stress simply because the rate of strain hardening is greater. Fundamental knowledge of the mechanisms of creep cannot, therefore, be obtained by noting the effect of the stress on the secondary creep rate, since for each stress a uniquely different structure is obtained.

The most characteristic feature of a thermally activated process is its activation enthalpy, h, where

$$g = h - Ts \tag{2}$$

and s is the activation entropy. Different mechanisms may occasionally have identical or almost equal activation enthalpies. But, often a clear identification of a single rate-controlling mechanism can be made in terms of the activation enthalpy. The approach to understanding creep through the activation enthalpy has indeed led to a considerably clearer picture of the processes that are involved. The activation enthalpy for creep can be deduced fairly accurately from easily obtainable data. As shown by Eqn. (1), a small and rapid change in temperature, holding the stress constant, will cause a change in strain rate such that

$$\frac{\Delta \ln \dot{\gamma}}{\Delta(-1/kT)} \simeq \frac{\partial \ln \dot{\gamma}}{\partial(-1/kT)} = \frac{\partial \ln f}{\partial(-1/kT)} + \frac{\partial(-g/kT)}{\partial(-1/kT)} \tag{3}$$

If the change in temperature is imposed sufficiently rapidly, the structure will not change. Furthermore, the frequency factor enters Eqn. (3) as a logarithmic term and is of itself only mildly, if at all, temperature dependent so that its contribution is negligible. Since the last term, according to the well-known relationship of classical thermodynamics, equals the enthalpy of activation,

the enthalpy of activation can be estimated from

$$h \simeq \frac{\Delta \ln \dot{\varepsilon}}{\Delta(-1/kT)} \tag{4}$$

A typical example of the type of data obtained in a change in temperature test[4] is shown in Fig. 3, the activation enthalpies that are recorded were calculated from Eqn. (4). For limited ranges of temperature and not too extensive variations in the creep rate, the activation enthalpies were generally found to be insensitive to the magnitude of the applied stress and the strain[4] as illustrated in Fig. 4. These observations immediately disqualify some previously held concepts that creep is a result of two parallel processes, the transient and the quasi-viscous, one of which, the transient, predominates in the primary stages of creep and the other of which, the quasi-viscous, predominates over the secondary stage. Neglecting the very small decrease in activation enthalpy with strain and the small insignificant trends with stress, the same activation enthalpy is obtained over both primary and secondary stages of creep. Thus, similar unit molecular processes are now thought to be responsible for the primary and secondary stages of creep. This can be rationalized by attributing the changes in creep rate to the effect of changes in structure. During the primary stage of creep the rate of strain hardening exceeds the rate of softening due to recovery whereas these rates are balanced over the secondary stage of creep. Thus, the primary stage "appears" to be transient and the secondary stage "appears" to be quasi-viscous.

As shown by the example given in Fig. 5, the activation enthalpies for creep often depend on the temperature of test.[5] Above about one-half of its melting temperature, the activation enthalpy for creep of high purity Al has a constant value of about 35,500 cal/mole. It also has approximately a constant value of about 28,000 cal/mole from about 240°K and 380°K. But between 380° and 500°K the activation enthalpy increases from 28,000 to 35,000 cal/mole

and from 0° to 240°K it increases from zero to about 28,000 cal/mole. Obviously, different mechanisms of creep operate over different ranges of temperatures. Over the ranges where the activation enthalpies are substantially constant, we must conclude, as we shall see later, that the enthalpy of activation is not only insensitive to the temperature but also to the stress and structure. There are at least two different ways in which increases in activation enthalpies can arise with increases in temperature. The first way[5] occurs when the total creep rate is the sum of two or more independent parallel processes each having different activation enthalpies. Over the lower temperature range, therefore, the creep rate is due primarily to the lower activation enthalpy process. At higher temperatures, however, successively higher percentages of the total creep rate are attributable to the higher activation enthalpy process. Thus, the average activation enthalpy increases with an increase in temperature. The second method in which the activation enthalpy can increase with increasing temperatures for even a single process is illustrated by the dislocation intersection mechanism.[6, 7] Here the enthalpy of activation is strongly dependent on the stress and structure. If, as shown by Eqn. (1), a constant strain rate is maintained, the activation enthalpy that need be supplied increases almost linearly with the absolute temperature according to

$$h = sT + kT \ln \frac{f}{\dot{\gamma}} \tag{5}$$

At first it may appear somewhat strange that an enthalpy of activation for a single mechanism can increase almost linearly with temperature. But the issue is easily resolved when it is recognized that the enthalpy is a function of the applied stress. What really happens is that the applied stress adjusts itself to achieve the imposed strain rate. For a given stress, the probability of a successful thermal fluctuation is decreased as the temperature decreases. To maintain the same strain rate at a lower temperature, therefore, a lower activation enthalpy is required. This is achieved by the appropriate increase

Table I

Tentative Identification of Creep Mechanisms in High Purity Aluminum

Region	Range °K	Observed Trend in H	Possible Mechanisms	References
A	0 to 150	H = const T	Intersection	6, 7, 8
B	150 to 240	More rapid than linear with T	Intersection and Cross-Slip	6
C	240 to 380	28,000 cal/mole	Cross-slip	8 to 13
D	380 to 500	Transition from 28,000 to 35,500 cal/mole	(a) Cross-slip and climb, or (b) Motion of jogged screw dislocations	
E	500 to 920 intermediate stresses	35,500 cal/mole	(a) Dislocation climb (b) Motion of jogged screw	
F	920 low stresses	35,500 cal/mole	Motion of jogged screws	

in stress to maintain an enthalpy that varies linearly with the temperature.

The data recorded in Fig. 5 can be described in terms of the ranges suggested in Table I. In this lecture major emphasis will be given to creep in Regions E and F where the temperature is above about one-half of the melting temperature. Nevertheless, a few remarks about the lower temperature regions appear to be appropriate here. Although there yet remains some questions relative to the details of strain hardening in Region A, there is strong evidence that the mechanism that controls the strain rate is the thermally activated intersection of dislocations. Extensive cross-slip is observed in Region C. Schoeck and Seeger[9] estimate theoretically that the activation energy for cross-slip in Al is about 24,000 cal/mole whereas Friedel[10] suggests that it is about 28,000 cal/mole the value being insensitive to the applied stress. In the same temperature range Aestrom[11], using the Borelius microcalorimeter, found that cold-worked Al will recover by a process that has an activation energy of 28,000 cal/mole. And recently, Jaffe and Dorn[12, 13] have presented additional evidence in favor of the suggestion that creep of Al in Range C is controlled by the cross-slip mechanism. Accordingly, Region B must then be ascribed to action of both the intersection and the cross-slip mechanisms. Regions E and F which appear to be characterized by a unique activation energy of 35,500 cal/mole will be discussed in detail from a theoretical viewpoint later, as will also the results in Region D.

Extensive investigations on the activation enthalpies for creep of a series of pure metals above one-half of their melting temperature have resulted in the correlation of this quantity with the activation enthalpies for self-diffusion as shown in Fig. 6.[2, 14, 15] (Formerly reported Al and Mg datum points have been omitted because the previously used ΔH_D for Al was an estimated value and because recent evidence[16] on the ΔH_c of Mg suggests the former correlation was fortuitous). Only one well-documented and reliable exception to the good agreement between the activation enthalpy for high temperature creep

and that for self-diffusion is now known, namely that for Mg.[16] A more extensive correlation[2, 14] of the activation enthalpies for creep, stress rupture and self-diffusion with the melting temperature of pure metals was recently extended by Bechtold, Wessel and France[17] as shown in Fig. 7. Although there may be some doubt regarding the accuracy of some of these data, particularly those deduced from stress-rupture tests, there is little doubt that, for most metals, the activation enthalpy for high temperature creep is indeed in close agreement with that for self-diffusion.

Numerous studies have also been made on the effect of stress on the creep rate and on the activation volumes[3, 7, 8, 13, 14, 18 to 20, 22] that are significant. Former investigations on the substructural changes during creep have been summarized by Sully[21] and currently there is much activity to uncover the significant substructural characteristics pertinent to creep with the aid of electron microscopy. These issues, however, are best understood in terms of theoretical models of creep mechanisms. And from this viewpoint the most significant feature is that the model for high temperature creep must be developed so as to give an activation enthalpy about equal to that for self-diffusion.

2. Vacancies and Diffusion

Over the years many theories for creep have been proposed. But, as shown in the preceeding section of this lecture, only those theories for high temperature creep that require enthalpies of activation which approximate those for self-diffusion can possibly be realistic. A brief review of diffusion will be given here in order to provide the basis for describing such theories.

Diffusion can take place by several mechanisms. As documented in a number of recent reviews,[23-25] on this subject, however, the volume diffusion in pure metals, and substitutional alloys as well, is now known to occur principally by the vacancy mechanism as shown in Fig. 8. The rate at which the radioactive atoms, shown by *, move is directly related to the number of vacancies that are present. The equilibrium number of single vacancies in a

crystal are easily deduced from statistical thermodynamics.[26-28] Since the equilibrium number of vacancies are independent of the mechanism whereby they are produced, it is permissible to view their production as given in the transition from "a" to "b" of Fig. 8. When an atom is removed from the near center of a crystal, the bonds with the adjacent coordinated atoms must be broken and when that atom is placed on the surface one-half of these bonds are restored. The total work involved in making a vacancy, therefore, is the work u_f of breaking one-half of the atomic bonds plus the work $p\Omega$ done against the surroundings where p is the pressure and Ω is the volume expansion per vacancy produced, a quantity almost equal to the atomic volume. The atoms adjacent to the vacancy now have a lower frequency of vibration than before. Consequently, the vibrational entropy changes an amount s_f per vacancy produced and the work that must be done to produce a single vacancy at a given site is

$$g_f = u_f + p\Omega_f - Ts_f \tag{6}$$

On the other hand, the increase in free energy upon introduction of n independent and non-interacting vacancies at random among n_A atoms, is

$$G_f = ng_f - kT \ln \frac{(n_A + n)!}{n_A!\, n!} \tag{7}$$

where the last term arises from the configurational entropy of random mixing of n vacancies with n_A atoms. Therefore, the average work that must be done in adding a vacancy to a random mixture of n vacancies and n_A atoms is

$$\frac{dG_f}{dn} = g_f - kT \ln \frac{n + n_A}{n} \tag{8}$$

Since the free energy is a minimum at equilibrium, the equilibrium number of vacancies, n_0, in a solid is given by $\left(\partial G_f/\partial n\right)_{n=n_0} = 0$

and consequently

$$n_0/n_s = e^{-g_f/kT} \tag{9}$$

where $n_s = n_A + n$ the total number of lattice sites.

The extra work that must be done in creating a vacancy in a crystal that contains n vacancies as compared with one that contains the equilibrium number of vacancies per unit volume is given by the chemical work term,

$$W_c = \left(\frac{\partial G}{\partial n}\right)_n - \left(\frac{\partial G}{\partial n}\right)_{n_0} = kT \ln \frac{n}{n_0} \tag{10}$$

This is, therefore, also the constant pressure isothermal work that a system which is supersaturated with vacancies can do.

Diffusion is a random walk phenomenon and, in the simple case of self-diffusion, it can be described in terms of the mass migration of tagged, radioactive, atoms. When the concentration of such tagged atoms is uniform, their random walk causes no net mass transfer. But when the situation is as shown in Fig. 9, a mass migration will take place in the opposite direction to the concentration gradient. Here N* is the number of tagged atoms per cm^2 of a crystal plane, where λ is the jump distance. Since all atoms are chemically identical each atom jumps with the same frequency f in any one fixed direction. Therefore, in time δt the increase in the number of tagged atoms on the plane at x is given by

$$\delta N^* = \{N^*(x-\lambda) + N^*(x+\lambda) - 2N^*\} f \delta t$$

or

$$\frac{\delta N^*}{\delta t} = D\left(\frac{\partial^2 N}{\partial x^2}\right)$$

where

$$D = f\lambda^2 \tag{11}$$

is the diffusivity as defined by Ficks second law for diffusion.

The frequency f depends on the mechanism of diffusion that takes place. For the vacancy mechanism of diffusion the frequency is given by the frequency ν of an atom adjacent to a vacancy, times the probability that it has enough energy to move into a vacant site, namely

$$e^{-g_m/kT} = e^{-\frac{U_m + p\Omega_m - S_m T}{kT}}$$

times the probability that the atom is moving in a given direction, α, times the probability, $\frac{n}{n_s}$, that a vacancy is in that direction. Therefore,

$$D = \alpha\lambda^2\nu\frac{n}{n_s}e^{-\frac{U_m + p\Omega_m - S_m T}{kT}} \tag{12}$$

For the simple cubic lattice the factor α, the probability of motion in any one of the six possible directions, is 1/6. The same type of analysis, with $n/n_s = 1$ reveals that the diffusivity of a vacancy is

$$D_v = \alpha\lambda^2\nu e^{-\frac{U_m + p\Omega_m - S_m T}{kT}} \tag{13}$$

When the diffusivity is determined under conditions where the concentration of vacancies is given by thermal equilibrium, the value n/n_s in Eqn. (12) must be replaced by n_o/n_s of Eqn. (9). Therefore,

$$D_o = \alpha \lambda^2 \nu e^{\frac{s_f + s_m}{k}} e^{-\left\{\frac{u_f + u_m + p(\Omega_f + \Omega_m)}{kT}\right\}}$$

$$= \alpha \lambda^2 \nu e^{s_d} e^{-\frac{(u_d + p\Omega_d)}{kT}} \qquad (14)$$

where the subscript d refers to the values for diffusion. Consequently, the activation enthalpy for diffusion, $u_d + p\Omega_d$ is equal to the sums of the enthalpies of formation and motion of vacancies. And the fact that the activation enthalpies for high temperature creep closely approximate those for self-diffusion strongly suggests that this creep is dictated by atomistic mechanisms that involve the formation and migration of vacancies.

Several different theories for high temperature creep that are based on the formation and migration of vacancies have been postulated thus far. The first, originally conceived by Nabarro[29] and elaborated upon by Herring[30] is known as stress directed diffusion of vacancies. It is applicable at very high temperatures, those approaching the melting temperatures, and at very low stresses, presumably below those which are required to operate a Frank-Read dislocation source or to generate dislocations by other means. The theory is quite accurately and completely formulated. It concerns the formation of vacancies under an applied stress at the grain boundaries normal to the applied stress and migration of the vacancies through the grain volume to the boundaries parallel to the applied stress. Atoms, of course, migrate in the opposite direction to the vacancies and thus provide the permanent strain. Several investigations on creep near the melting temperature give creep rates that are in good agreement with the theory. In view of the fact that this mechanism of creep is now well established, it will not be presented in detail here.

Whereas creep by stress directed diffusion of vacancies can occur in amorphous as well as crystalline solids, since dislocations do not participate in this mechanism, the remaining two theories, (a) thermally activated motion of jogged screw dislocations, and (b) the climb of edge dislocations are limited to crystalline materials since they both depend on dislocation processes. The validity of these theories is not nearly so well established as that for the stress directed diffusion of vacancies and consequently they are yet under discussion. They will be reviewed in the next sections of this lecture so as to provide a basis for discussing the problems that yet restrain their unqualified acceptance.

3. Creep due to the Motion of Jogged Screw Dislocations

The thermally activated motion of jogged screw dislocations has been described by Mott,[31] Seeger,[32] Friedel[33] and van Bueren.[34] The following analysis is taken, with minor modifications, principally from the recent investigations by Hirsch and Warrington.[35]

As shown in Fig. 10, screw dislocations bow out under an average local stress $\tau - \tau^*$ where τ is the applied stress and τ^* is a long range back stress, the dislocation being held up at the jogs. At "a" is shown a unit vacancy-forming jog whereas at "b" a unit interstitial-forming jog is shown, In view of the high energy, however, of forming interstitials, the operative process of creep must depend primarily on the vacancy mechanism. At both "a" and "d" a force $(\tau - \tau^*)\ell_j b$ is operative where ℓ_j is the mean distance between the jogs. If the jog at "a" is to move forward one Burgers vector, an atom must jump into position V_1 and form a vacancy in the lattice. Consequently, point "a" is a source of vacancies. In order for the jog at "d" to move forward, a lattice vacancy must exchange with the atom at B_1. Consequently, point "d" is a vacancy sink. Throughout the lattice there are many sources and sinks. For example, dislocations and grain boundaries can also serve as sources and

sinks for vacancies.

Since it is the objective here to calculate the component of the creep rate due exclusively to the motion of jogged screw dislocations it appears appropriate to neglect all other non-conservative vacancy mechanisms for the present.

The question now arises as to how vacancies generated at sources diffuse to the nearby sinks. For example, Lothe[36] recently suggested that inasmuch as the energy to form a vacancy on a dislocation is smaller than that to form a vacancy some distance from a dislocation, and inasmuch as pipe diffusion along the dislocation is much more rapid than volume diffusion, pipe diffusion must control the process. Although such a mechanism undoubtedly contributes to the overall process, there are several factors that suggest it may not be important. First, pipe diffusion would lead to an activation enthalpy for creep that is considerably below that for volume self-diffusion. Secondly, pipe diffusion must take place along the single atomic path coincident with the dislocation line whereas volume diffusion may proceed along many paths, a factor which compensates for the higher activation enthalpy of volume diffusion. Furthermore, the osmotic work that can be done by supersaturation in the vicinity of a vacancy forming jog may be sufficient to cause large numbers of vacancies to leave the dislocation pipe to undertake volume diffusion.

For superjogs having a height of p atomic planes, the average force acting on each atomic height of the jog is $(\tau-\tau^*)\ell_j b/p$. The energy that must be supplied by a thermal fluctuation in order to form a vacancy is therefore $g_f-(\tau-\tau^*)\ell_j b^2/p$ where g_f is the free energy of the formation of a vacancy and the last terms represents the work contributed mechanically by the force acting at the jog. But to provide a permanent forward motion of the jog, the vacancy so produced will, according to the present analysis, have to move away. Consequently, the frequency ν_{vj}^{+} with which a vacancy forming superjog moves forward one

Burgers vector is given by

$$\nu_{vj}^{+} = \nu e^{-\frac{\{g_f - (\tau - \tau^*) l_j b^2/p\}}{kT}} Z e^{-\frac{g_m}{kT}} \tag{15}$$

where ν is only slightly less than the Debye frequency and Z is the coordination number. Having so jumped forward, the jog may now jump back again. This frequency depends on the probability that a vacancy is next to the last atom on the jog times the probability that it will exchange with that atom and thus return the jog to its original position. Consequently, the frequency for the reverse reaction is

$$\nu_{vj}^{-} = \nu Z \frac{n^+}{n_s} e^{-\frac{g_m}{kT}} \tag{16}$$

where $\frac{n^+}{n_s}$ is the actual probability of finding a vacancy adjacent to a vacancy forming jog. But

$$\frac{n^+}{n_s} = \frac{n^+}{n_0}\frac{n_0}{n_s} = \frac{n^+}{n_s} e^{-\frac{g_f}{kT}} \tag{17}$$

And, therefore, the net frequency of the forward motion of a vacancy forming jog, $\nu_{vj} = \nu_{vj}^{+} - \nu_{vj}^{-}$ is

$$\nu_{vj} = \nu Z e^{-\frac{g_d}{kT}} \left\{ e^{\frac{(\tau - \tau^*) l_j b^2/p}{kT}} - \frac{n^+}{n_0} \right\} \tag{18}$$

when the concentration of vacancies reaches the equilibrium concentration at a jog under a stress

$$(\tau - \tau^*) l_j b^2/p = kT \ln \frac{n^+}{n_0} \tag{19}$$

and the net frequency of the forward motion of a vacancy forming jog becomes zero. As a result of diffusion, however, n^+/n_0 never reaches this equilibrium value.

The forward frequency for motion of an interstitial forming jog can be obtained by the same method of analysis as that for a vacancy forming jog. The probability of finding a vacancy adjacent to an interstitial forming jog is given by

$$\frac{n^-}{n_s} = \frac{n^-}{n_0} \frac{n_0}{n_s} = \frac{n^-}{n_0} e^{-\frac{g_f}{kT}} \tag{20}$$

where $n^- < n^+$ since the interstitial jog is a vacancy sink. The net frequency of the forward motion of an interstitial forming jog, therefore, is given by

$$\nu_{ij} = \nu Z e^{-\frac{g_d}{kT}} \left\{ \frac{n^-}{n_0} e^{\frac{(\tau - \tau^*) l_j b^2/p}{kT}} - 1 \right\} \tag{21}$$

Although it might be expected that the interstitial forming jogs might lag behind the vacancy forming jogs, this does not occur since as one lags behind the other, the force due to the line tension of the dislocation segment increases. Consequently, jogs of the same size can move forward with the same net frequency.

Superjogs having the greater heights p will lag behind those having the smaller heights because the force causing the production or adsorption of vacancies is less per vacancy. When p is not too large, however, the same adjustment of forces on each jog will take place as previously discussed. Very large superjogs, however, will not be able to keep up with the advancing screw dislocation and may therefore form dipole pairs of positive and negative edge dislocations, particularly once the applied stress is removed. The rate of removal of such dipoles will depend on the climb of edge dislocations which will be described in the following section of this lecture.

Consequently, for small jogs either Eqn. (18) or (21) gives a good estimate of the frequency of the forward motion of the jog. The strain rate, $\dot{\gamma}$, is given by

$$\dot{\gamma} = \left(\frac{\rho_s}{\ell_j}\right)(\ell_j b)\, b\, \nu_{vj} = \rho_s\, b^2 \nu_{vj} \tag{22}$$

where ρ_s is the total length of screw dislocations per cm^3, $\frac{1}{\ell_j}$ is the number of jogs per cm^3 on screw dislocations and ρ_s/ℓ_j is the area swept out per activation. Introducing the value of ν_{vj} given by Eqn. (18) reveals that

$$\dot{\gamma} = \rho_s b^2 \nu Z e^{-\frac{g_d}{kT}} \left\{ e^{\frac{(\tau-\tau^*)\ell_j b^2/p}{kT}} - \frac{n^+}{n_0} \right\} \tag{23}$$

When the applied stress is small, the rate of production of excess vacancies can be expected to be much smaller than the rate at which they diffuse to sinks. Under these conditions $\frac{n^+}{n_0}$ approaches unity. At higher strain rates, however, $\frac{n^+}{n_0}$ increases. A satisfactory analysis of the dependency of $\frac{n^+}{n_0}$ on the frequency of motion of a jog has not yet been made. Hirsch and Warrington,[35] however, suggest that $\frac{n^+}{n_0}$ never becomes very great. Since $\frac{n^+}{n_0}$ should increase with the applied stress according to

$$\frac{n^+}{n_0} \propto e^{\frac{(\tau-\tau^*)\ell_j b^2/p}{kT}}$$

Eqn. (23) can be written as

$$\dot{\gamma} = \frac{\rho_s b^2 z \nu}{F} e^{-\frac{g_d}{kT}} e^{\frac{(\tau-\tau^*)\ell_j b^2/p}{kT}} \tag{24}$$

where

$$\frac{1}{F} = 1 - \frac{n^+}{n_0} e^{-\frac{(\tau-\tau^*)\ell_j b^2/p}{kT}} \tag{25}$$

and F is never very large.

Under equilibrium conditions the probability of finding a jog is given by

$$b/\ell_j = e^{-g_j/kT} \tag{26}$$

where g_j is the free energy of formation of a jog. But, as the screw dislocations sweep through the crystal and intersect forest screw dislo-

cations, additional jogs are produced. Furthermore, additional superjogs are also produced by cross-slip. In the absence of an applied stress such excess jogs could combine or migrate to nodal points so as to preserve the validity of Eqn. (26). But under the action of a stress such migration and annihilation of jogs becomes difficult. Consequently, after some high temperature deformation Eqn. (26) may no longer be valid and ℓ_j may be smaller than the true equilibrium value. For this model therefore, ℓ_j represents one of the structure sensitive parameters of creep.

It is difficult to prescribe a priori whether long range back stresses will be significant or not. In cases where the stacking fault energy is high or intermediate, the screw dislocations will cross-slip with relative ease at the high temperatures that are being considered here. Screw dislocations may also enter twist boundaries which have only local stress fields and, therefore, may not contribute significantly to τ^*. For these reasons it appears appropriate to neglect any small effect of long range back stresses and suggest that

$$\dot{\gamma} = \frac{\rho_s b^2 \nu Z}{F} e^{\frac{\tau \ell_j b^2/p}{kT}} \tag{27}$$

where ℓ_j depends on the average distance each screw dislocation has moved. The change in creep rate during the course of creep then must depend on the changes in the two structure sensitive parameters ρ_s and ℓ_j/b. It can be expected that during the course of creep under constant stress the edge dislocations become blocked leading to the observed formation of subgrains. Consequently, ρ_s should decrease with continued creep. But according to this model, creep should decrease as the edge dislocations pile up and thus decrease ρ_s. It appears that the climb model must also be invoked in order to account for continued straining during secondary creep.

The most satisfactory confirmation of the jogged screw dislocation model for high temperature deformation is contained in the paper by Hirsch and Warrington. They contend that below about one-half of the melting temperature, diffusion is negligible, and the deformation occurs by an athermal process, namely, the formation of vacancies at jogs.

Hirsch has suggested, without furnishing complete proof, that interstitial forming jogs, in contrast to vacancy forming jogs, can slide conservatively parallel to the screw dislocation and thus become annihilated at nodal points. It appears, however, that such sliding will be quite difficult when a stress is applied that causes the dislocation to bow out. It is reasonable to believe, however, that the interstitial forming jogs may be left in the wake of a moving dislocation. In either event the flow stress will depend essentially on the work required to produce vacancies, as given by

$$\tau_a = \frac{g_f}{l_j b^2/p} \tag{28}$$

where again long range back stresses are assumed to be small. When $\tau l_j b^2/p > kT$ the flow stress for the thermally activated process, as shown by Eqn. (27) is

$$\tau = \frac{g_d}{l_j b^2/p} - \frac{kT}{l_j b^2/p} \ln \frac{P_s b^2 \nu Z}{\dot{\gamma}} \tag{29}$$

Since g_f is less than g_d by g_m the athermal process prevails at temperatures below a critical value T_c given by

$$g_m = k T_c \ln \frac{P_s b^2 \nu}{\dot{\gamma}} \tag{30}$$

These trends are shown schematically in Fig. 11. Formulating the Cottrell-Stokes ratio between the flow stress for a given state for various temperatures above T_c and the flow stress at T_c gives

$$\frac{\tau}{\tau_a} = 1 - \frac{k(T-T_c)}{g_f} \ln \frac{\rho_s b^2 \nu Z}{\dot{\gamma}} \tag{31a}$$

for $T > T_c$ and

$$\frac{\tau}{\tau_a} = 1 \tag{31b}$$

for $T < T_c$. Typical results are shown in Fig. 12 which confirm, with other details, the nominal validity of the model. As shown by Eqn. (27) the activation enthalpy for the higher stress levels is an amount $(\tau - \tau^*) l_j b/p$ less than that for self-diffusion. However, due to the rapid linear drop in stress with increase of temperature, the activation enthalpy rapidly approaches that for self-diffusion as the temperature increases when the strain rate is constant.

4. Creep as a Result of Climb of Edge Dislocations

Creep arising from the climb of edge dislocations is much more difficult to formulate theoretically than that arising from the motion of jogged screw dislocations. The major problem concerns the fact that the rate at which edge dislocations climb is determined by the intimate details of dislocation patterns produced under a stress. Since these patterns are not well known and since many different dislocation arrangements are possible, the following analysis can only represent the general trends.

The motion of jogged screw dislocations that was described in the preceeding section of this lecture is shown schematically in Fig. 13 as a plan view of the slip plane. When a stress is applied to a grain, dislocation sources such as that shown in "a" begin to operate. As indicated in "b" the edge dislocations

move rapidly whereas the jogged screw dislocations move much more slowly as described in the preceeding section of this lecture. But the edge dislocations may be arrested at various barriers. As this happens ρ_s the density of the active screw dislocations decreases.

It is well-known that during creep of most metals, subgrains are produced. Such subgrains must arise from the non-homogeneous distribution of stresses in each grain, resulting in localized bending and subsequent polygonization. Whereas the activation enthalpy for creep of unbent single crystals of Al remains at the value of 28,000 cal/mole for all temperatures above one-half of the melting temperature[37] the activation enthalpy for polycrystalline Al is about 35,500 cal/mole. Therefore, bending and polygonization play an important role in the high temperature creep of polycrystalline aggregates. Undoubtedly when the applied stress is increased the non-homogeneous local bending stresses also increase and produce the finer subgrain sizes that are observed. During pile up of the edge dislocation, the screw dislocations may enter twist boundaries or they may cross slip and be annihilated by joining screw dislocations with Burgers vectors of the opposite sign. During piling up of the edge dislocations, as shown schematically in "d" of Fig. 13, the length ρ_s per unit volume of the screw dislocations decreases and the creep rate diminishes, as is actually observed over the primary stage of creep. Although this model has certain merits, it does not appear to coincide exactly with the typical observations made with the electron microscope. The most prominent substructural feature of crept metals is the formation of subgrains; pile ups of edge dislocations are seldom seen excepting in metals having extremely low stacking fault energies. Nevertheless, something analogous to pile ups must form, since the applied stress will cause clusters of dislocations of the same sign to form in front of any substantial barrier. The approximation that such clusters can be represented by an idealized pile up will be made in order to facilitate the analysis.

When the edge dislocations climb they relieve the back stresses on the dislocation sources and permit, in this way, the generation of additional dislo-

cations. As will be described, it is expected that both the motion of jogged screw dislocations and the climb of edge dislocations contribute to the steady state creep observed over the range of the secondary stage.

A suggested model of the piled-up edge dislocations within a subgrain is shown in Fig. 14, where δ is the diameter of the subgrain. Since array "A" does not encounter any additional barriers within the subgrain, the edge dislocations move rapidly to the subgrain boundary. The boundary, however, is produced by the non-homogeneous bending stress fields in the total grain and consequently is a good sink for edge dislocations. Since edge dislocations enter the boundary easily, the contributions to creep of isolated sources of Type "A" must be ascribed to the motion of jogged screw dislocations. Consequently, ρ_s now refers to the density of screw dislocations issuing from sources of Type "A".

Dislocations from all other types of arrays, however, are held up at barriers and cannot, therefore, continue to the sub-boundary without first undertaking climb. The most likely general type of barrier results from the mutual interaction stress fields of the dislocations themselves. Only those arrays having the greatest number of dislocations will climb rapidly. Consequently, arrays of Type "B" will contribute much more to creep by climbing than those of Type "C", etc.

It is, therefore, reasonable to estimate the contribution to creep by climb as resulting principally from arrays of Type "B". The actual creep rate, however, will be somewhat larger than the value estimated in this way as a result of contributions of other arrays such as Type "C", etc.

The climb of edge dislocations has been described in a number of research papers.[38, 39] But the following discussion is based on an extension of Weertman's[40] method of analysis with the added details concerning the effect of structure on the process.

The velocity with which a dislocation climbs is given by

$$\dot{y} = b\nu_c \tag{32}$$

where ν_c is the frequency with which each segment of the edge dislocation climbs one Burgers vector. If the distance between the arrays of Type "B" is h, each dislocation must climb a height h/2. Consequently, the time it will take for a single dislocation to climb the required height is

$$t_c = \frac{h}{2b\nu_c} \tag{33}$$

If N is the number of arrays of type "B" per subgrain, the area swept out per unit volume is δ^2/δ^3 and the creep rate due to climb is given by

$$\dot{\gamma}_c = \frac{Nb}{\delta t_c} = \frac{2Nb^2\nu_c}{\delta h} \tag{34}$$

Whereas δ, as previously described, decreases as the stress increases, N increases with the stress and also with δ.

The statistical analyses for determining how N and δ depend on the stress have not yet been made. Nevertheless, the dependency of h on the stress is easily estimated in terms of well-known methods of analysis. Stroh[40] has shown that a stress concentration factor equal to $\sqrt{L/l}$ is obtained near the spur of a dislocation array, where L is the length of the array and l is the distance from the array at which the local stress is desired. Taking L as $\delta/4$ and l as h, the local stress barrier of the lower array that a moving dislocation of the upper array must surmount is given by $\tau\sqrt{\delta/4h}$. But, as shown by Eshelby, Frank and Nabarro[42] the stress acting on the leading dislocation of the upper array is equal to n times the applied stress where n is the number

of dislocations in the array of length $\delta/4$ which is given by

$$m = \pi \delta \tau (1-\mu)/4Gb,$$

where μ is Poisson's ratio.

Therefore, when

$$\frac{\pi \delta \tau^2 (1-\mu)}{4Gb} < \tau \sqrt{\frac{\delta}{4h_m}} \tag{35}$$

the leading dislocations of each array will be arrested. To simplify the analysis it will be assumed that the average height of climb is

where

$$h = \frac{2G^2 b^2}{\pi^2 (1-\mu)^2 \delta \tau^2} \tag{36}$$

And, therefore, the creep rate due to climb is given by

$$\dot{\gamma}_c = \pi^2 (1-\mu)^2 N \left(\tau/G\right)^2 \nu_c \tag{37}$$

The frequency with which climb takes place, ν_c, can now be estimated using almost the same procedures as those that have already been used in the preceeding section of this lecture for the motion of a jogged screw dislocation. The major difference between climb and the motion of jogged screw dislocation is revealed in Fig. 15. Whereas the motion of a jogged screw dislocation is stimulated by the shear stress, the climb of an edge dislocation is influenced only by the stress σ_{xx} nomal to the edge dislocation half plane as shown. When downward climb, designated +, takes place, a vacancy must be produced at a jog and moved away. The stress σ_{xx} during this process assists by doing the work $\sigma_{xx}\Omega_f$ where Ω_f is almost equal to the atomic volume. Therefore, the frequency of the downward (+) climb

is given by

$$\nu^{+} = \nu e^{-\frac{g_f - \sigma_{xx}\Omega_f}{kT}} z e^{-\frac{g_m}{kT}} p_j$$

where p_j is the probability of finding a jog. The reverse climb involves moving a vacancy into position at the jog and it is given by

$$\nu^{-} = \nu \frac{n^{+}}{n_s} z e^{-\frac{g_m}{kT}} p_j$$

Therefore, the net frequency for climb is

$$\nu_c = \nu z p_j e^{-\frac{g_f + g_m}{kT}} \left\{ e^{\frac{\sigma_{xx}\Omega}{kT}} - \frac{n^{+}}{n_s} \right\} \quad (38)$$

Two major problems arise concerning the details of Eqn. (39). The first concerns the stress σ_{xx} and the second the term p_j. The local stress σ_{xx} is sensitive to the structure and arises from the sums of the stress fields of the near dislocations. According to the simplified model that was assumed, here, it is given by two times the stress concentration factor times the stress field of a single edge dislocation. Therefore, to a fair approximation

$$\sigma_{xx} = 2\sqrt{\delta/4b} \left\{ \frac{Gb}{2\pi(1-\mu)h} \right\}$$

or, upon replacing the value of h by that given in Eqn. (36)

$$\sigma_{xx} = \frac{\pi^2 (1-\mu)^2 \delta^2 \zeta^3}{4\sqrt{2}\ G^2 b^2} \tag{39}$$

The analysis for $\frac{n^+}{n_0}$ might closely follow that given previously in the preceeding section of this lecture. In terms of the arguments presented there and in view of the approximations that have been made here it is not unreasonable to let n^+/n_0 be approximately unity.

The analysis for determining the probability of finding a jog on the edge dislocations has not yet been well formulated. Under conditions of forced climb the existing jogs are exhausted; therefore, new jogs must be produced if climb is to continue. In general, the nucleation of new jogs is quite difficult for low stacking fault energy crystals. As shown in Fig. 16 a high energy is required to add a single vacancy to a dislocation separated into its partials. If, however, two vacancies are added the maximum free energy is probably exceeded and a stable configuration of two jogs is produced. As additional vacancies add to the jogs, the partials of the dislocation segment between the jogs will continue to separate with a further decrease in energy. It is suggested that this may be the reason why the activation energy for recovery and re-crystallization of Al under zero stress is about twice that for self-diffusion.[43]

Under the action of stress concentrations due to piled-up arrays, however, the partial dislocations of the leading dislocations are brought more closely together. Under these conditions the thermal energy to produce a jog becomes smaller due to the reduction in the constriction energy. If the nucleation of jogs is not too difficult and if the rate of climb is not too rapid, the probability of finding a jog may approach the equilibrium value

$$P_j = b/\ell_j = e^{-g_i''/kT} \tag{40}$$

where g_i'' has a value of about 2,300 cal/mole, the jog energy in the absence of a constriction. To within the accuracy of these approximations, therefore, the creep rate for climb is given by

$$\dot{\gamma}_c = \pi^2 (1-\mu)^2 \nu Z N \left(\frac{\tau}{G}\right)^2 e^{-\frac{(g_d + g_j'')}{kT}} \left\{ e^{\frac{\pi^2 (1-\mu)^2 \delta^2 \tau^3 \Omega}{4\sqrt{2}\, G^2 b^2 kT}} - e^{-\frac{\pi^2 (1-\mu)^2 \delta^2 \tau^3 \Omega}{4\sqrt{2}\, G^2 b^2 kT}} \right\} \tag{41}$$

In general, the exponential term is small relative to unity and, therefore,

$$\dot{\gamma}_c = \frac{\pi^4 (1-\mu)^4 N \delta^2 Z \nu \tau^5 \Omega}{2\sqrt{2}\, b^2 G^4 kT} e^{-\frac{(g_d + g_j'')}{kT}} \tag{42}$$

This theory compares favorably with the experimental results. The activation enthalpy for creep is only slightly greater than that for self-diffusion. Furthermore, the effect of stress on the creep rate is also in good agreement with creep data. The results summarized recently by Sherby[44] as shown in Fig. 17 reveal that the creep rate indeed depends on about the fifth power of the stress.

Discussion

A rather complete theory is now available to describe the component of creep arising from the thermally activated motion of jogged screw dislocations in pure metals. And this theory appears to agree well with the experimental facts in the near neighborhood of about one-half of the melting temperature.

On the other hand, it is unlikely that creep occurs exclusively by this mechanism. As creep continues it is to be expected that some dislocation sources will become blocked as a result of the back-stress fields of interacting edge dislocations. Climb of such blocked dislocations can result in a second

component of creep. Consequently, it might be expected that both motion of jogged screw dislocations and the climb of edge dislocations contribute to the total observed creep. Whereas the activation enthalpy for the motion of jogged screw dislocations is slightly less than the activation enthalpy for diffusion, that for climb is slightly greater than that for self-diffusion. The mean value might, therefore, approximate that for self-diffusion as is observed experimentally.

Many problems concerning the effect of structural factors on the creep rate for climb must yet be studied in detail before a satisfactory theory can be formulated. A major problem concerns the nature and distribution of non-homogeneous stressing of each grain in the polycrystalline aggregate. This is undoubtedly the principal factor leading to subgrain formation. The size of such subgrains and the distribution of dislocation sources in the subgrains is also significant. Details of the effect of stress on the separation of partial dislocations and the resulting jog energies must also be known. Furthermore, the kinetics of nucleation of jogs and their diffusion demands additional attention. In conclusion, it is reasonable to state that only the first crude steps have been taken in an attempt to understand the high temperature creep of pure metals. But the progress looks good and some satisfactory solutions should soon become available.

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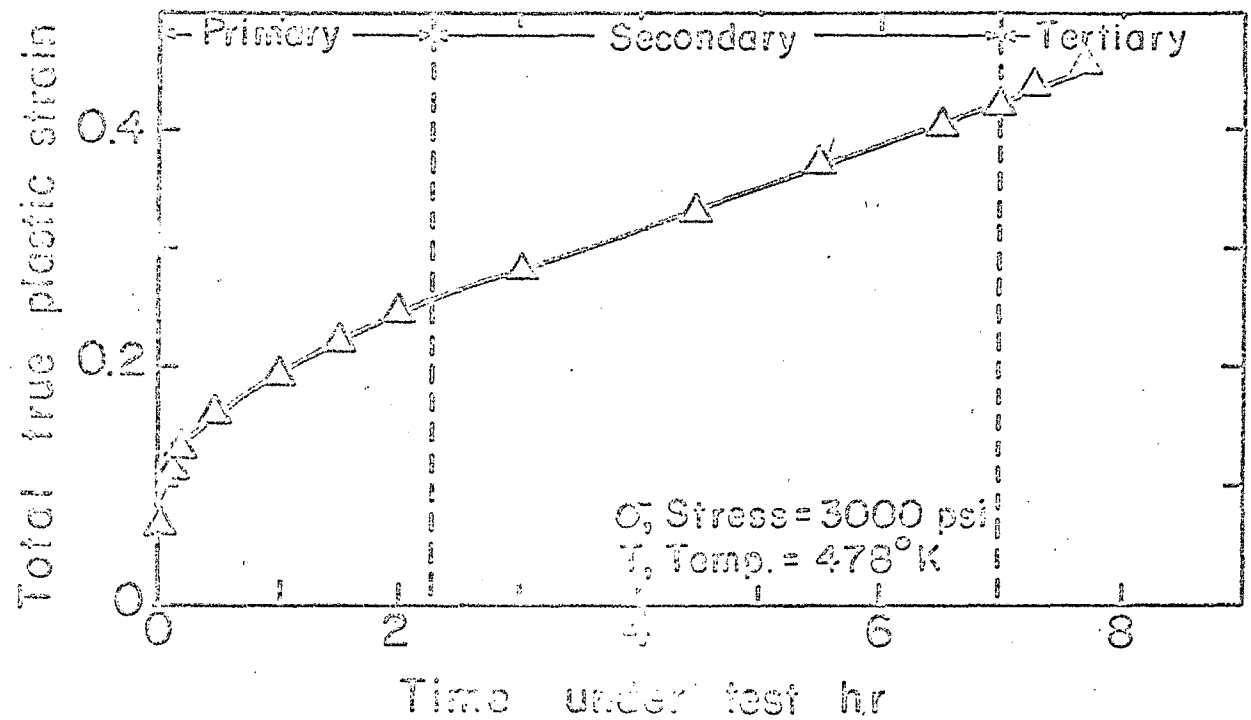


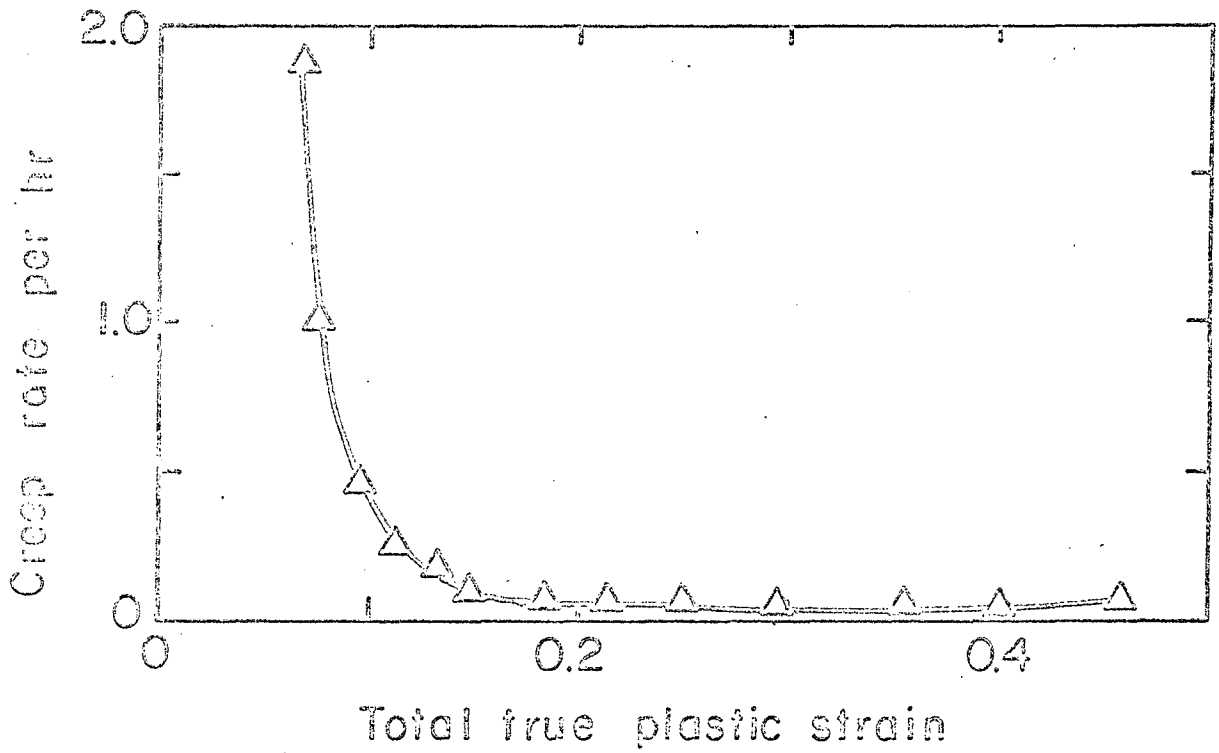


Fig. 1A & 1B Typical Creep Curve (high purity polycrystalline Al),[2] and Change in Creep Rate during Creep at Constant Stress and Temperature[2]

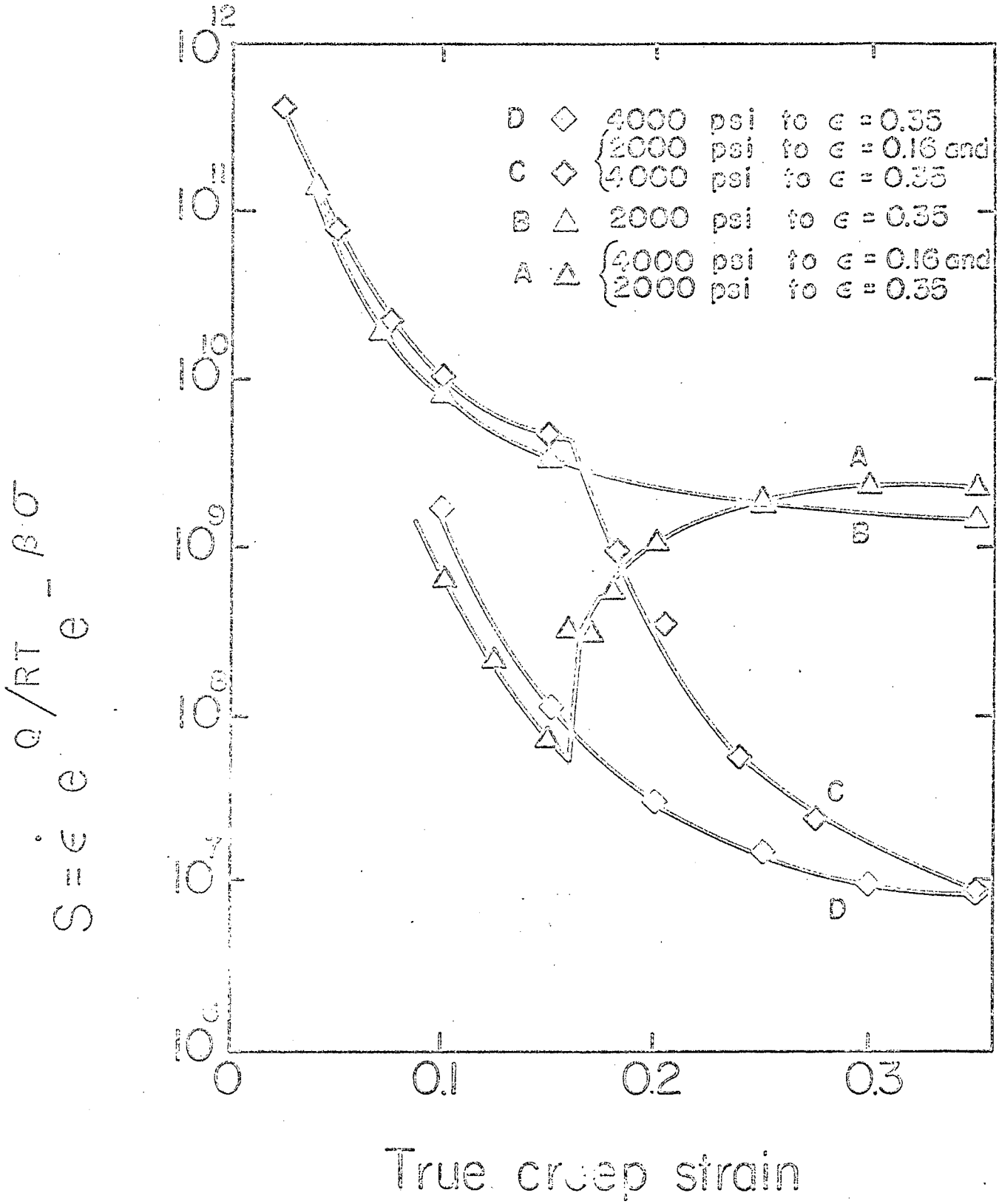
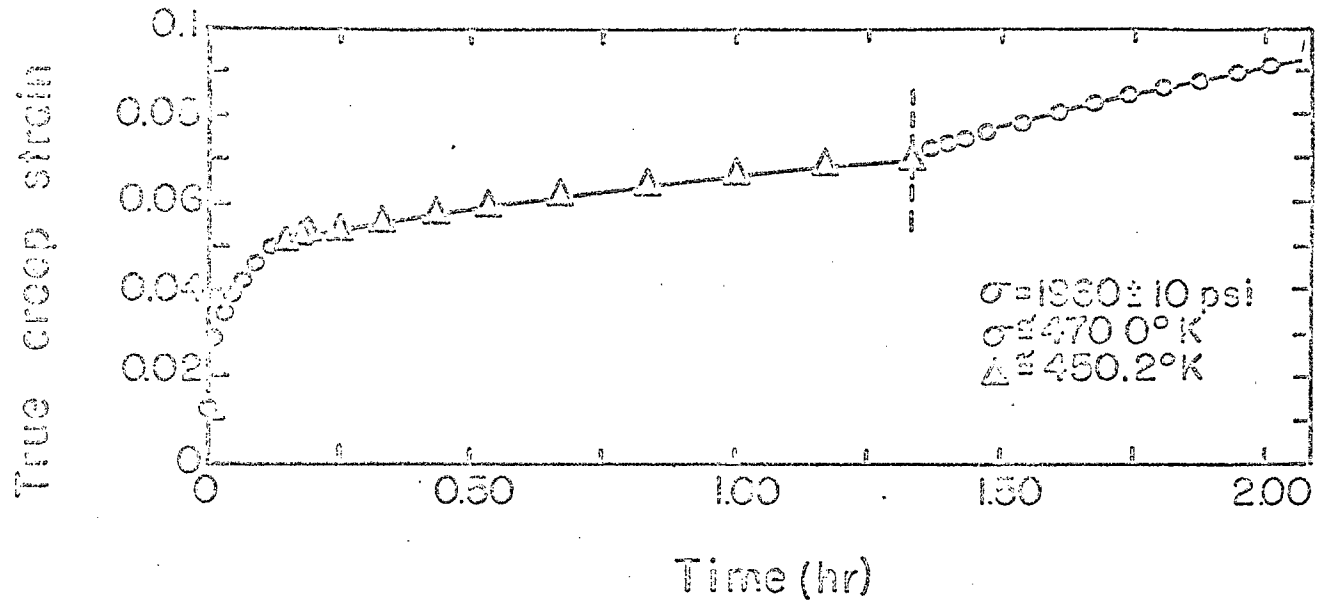


Fig. 2 Effect of Change of Stress on the Structure Parameter - True Creep Strain Relationship for High Purity Aluminum [3]



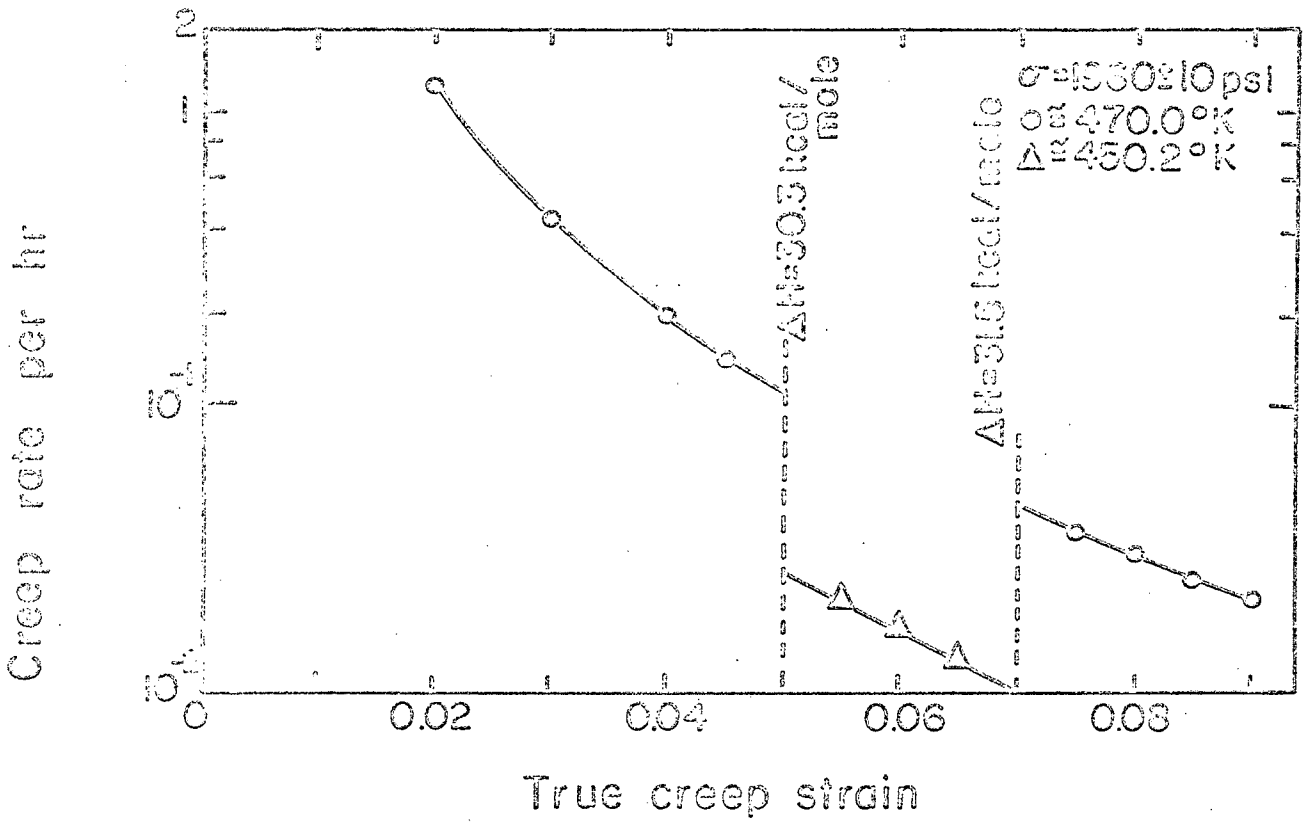


Fig. 3A & 3B Typical Creep Curve for Pure Aluminum under Cyclic-Temperature Condition,[4] and Creep Rate - True Creep Strain Curves for Use in the Evaluation of ΔH for Pure Aluminum (deduced from Fig. 3A above).[4]

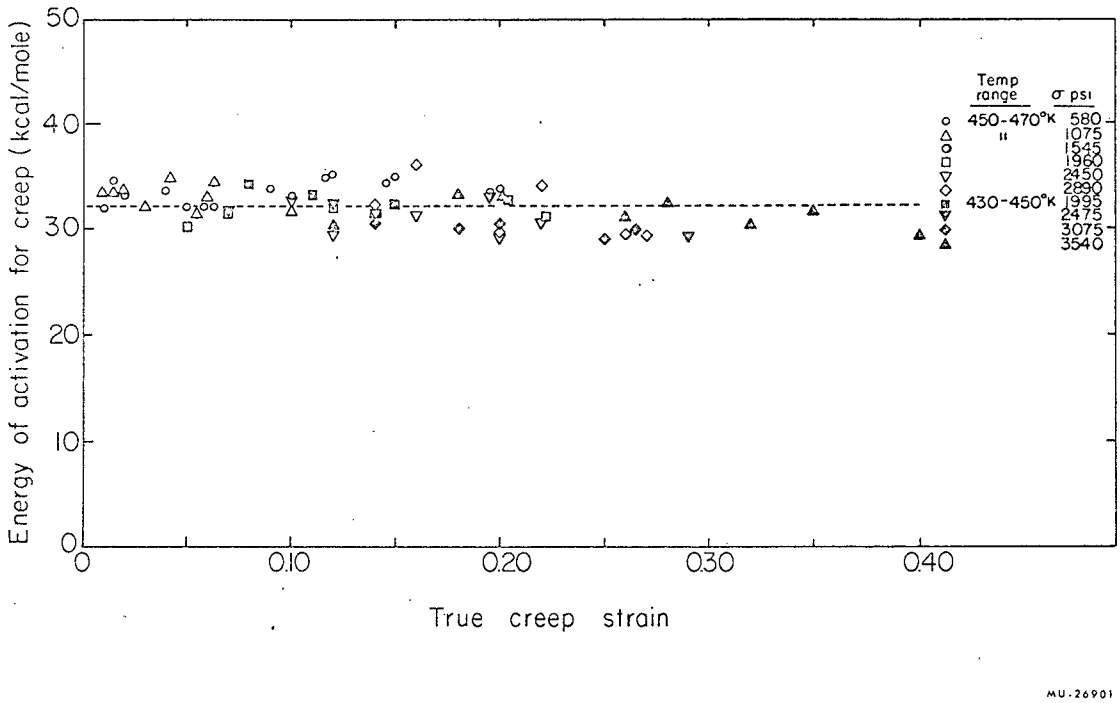


Fig. 4 Energy of Activation for Creep of Pure Aluminum as a Function of Creep Strain[4]

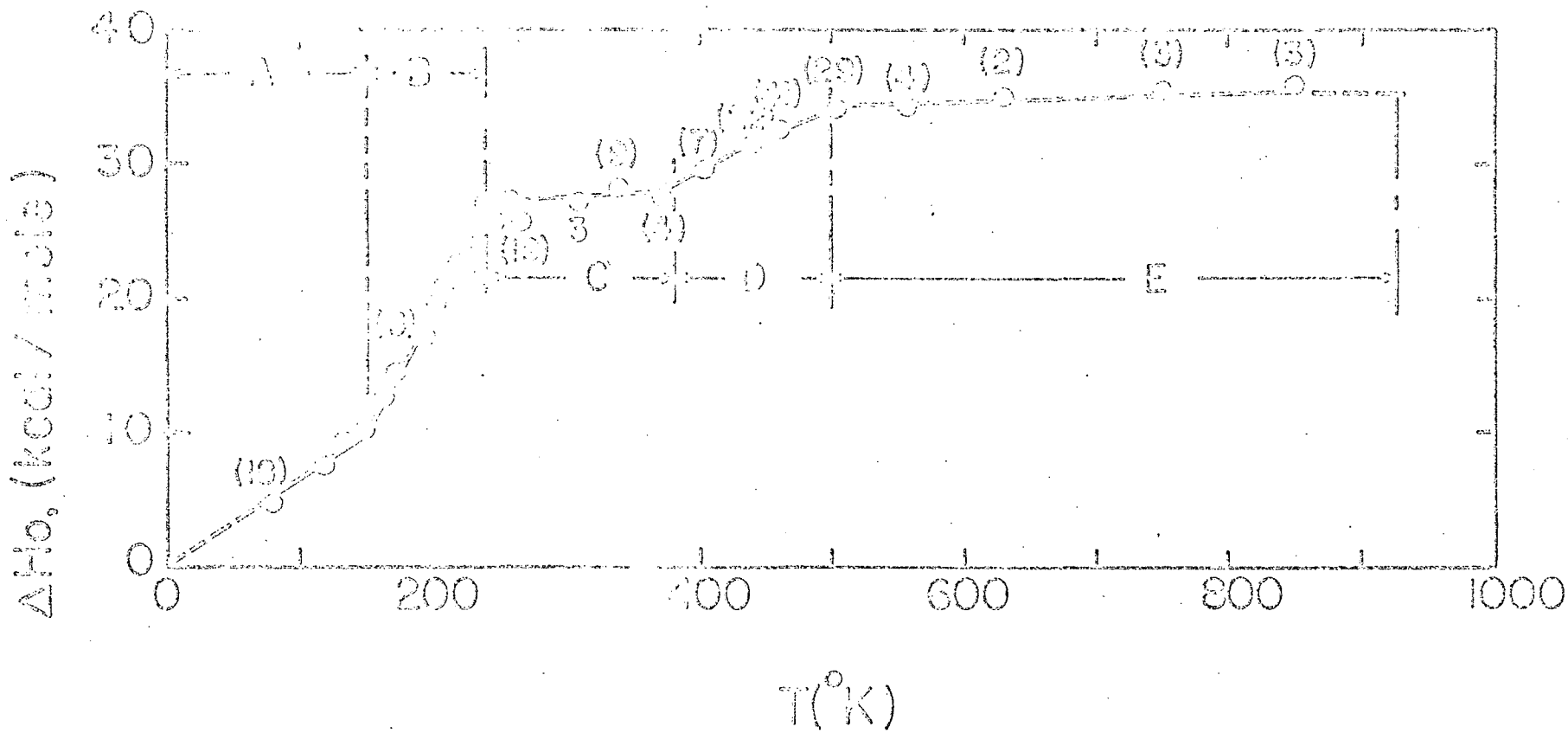


Fig. 5 Activation Energies for Creep of Pure Aluminum as a Function of the Absolute Temperature.[5]

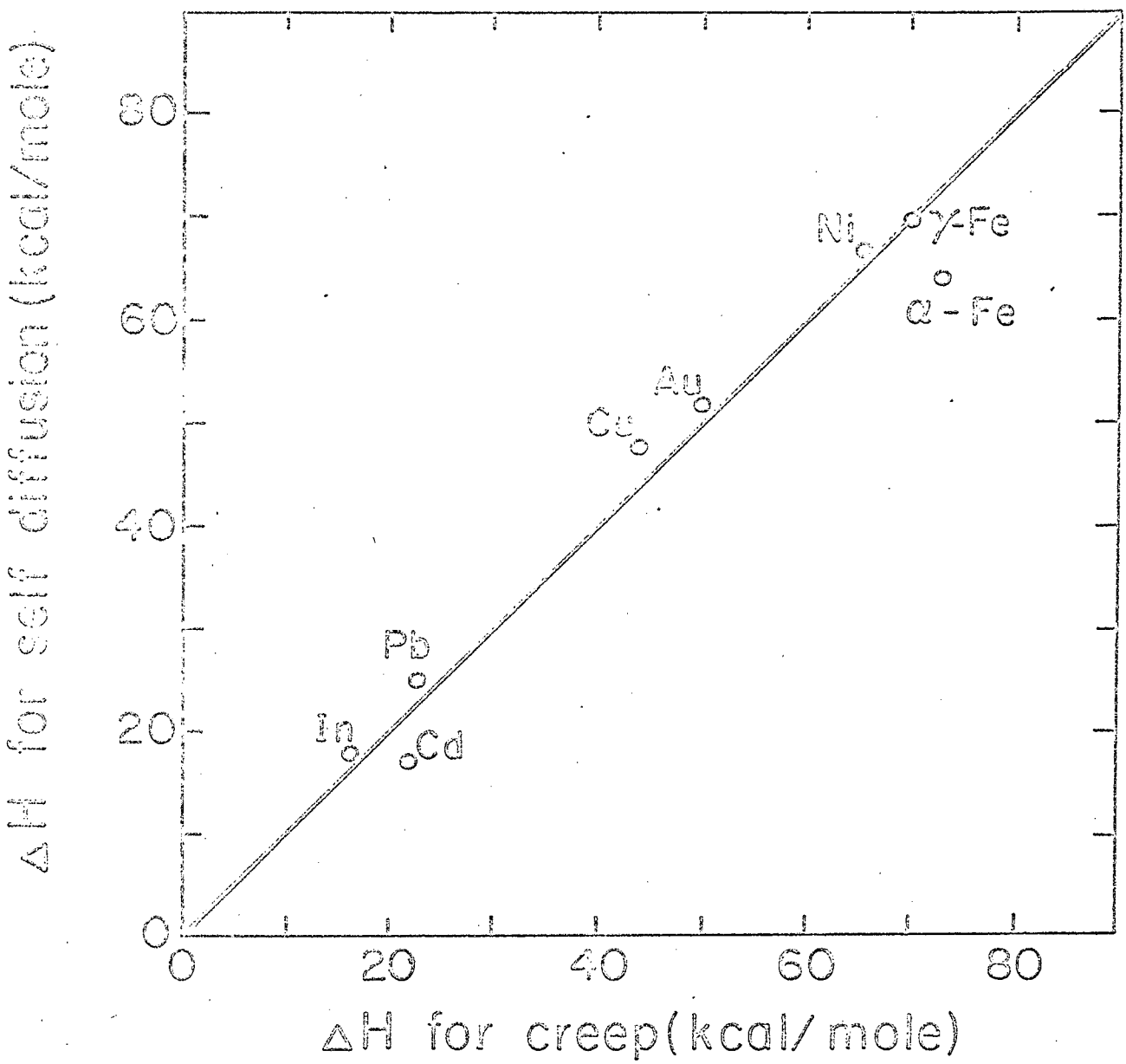


Fig. 6 Activation Energies for High Temperature Creep and Self-Diffusion.[2, 14, 15]

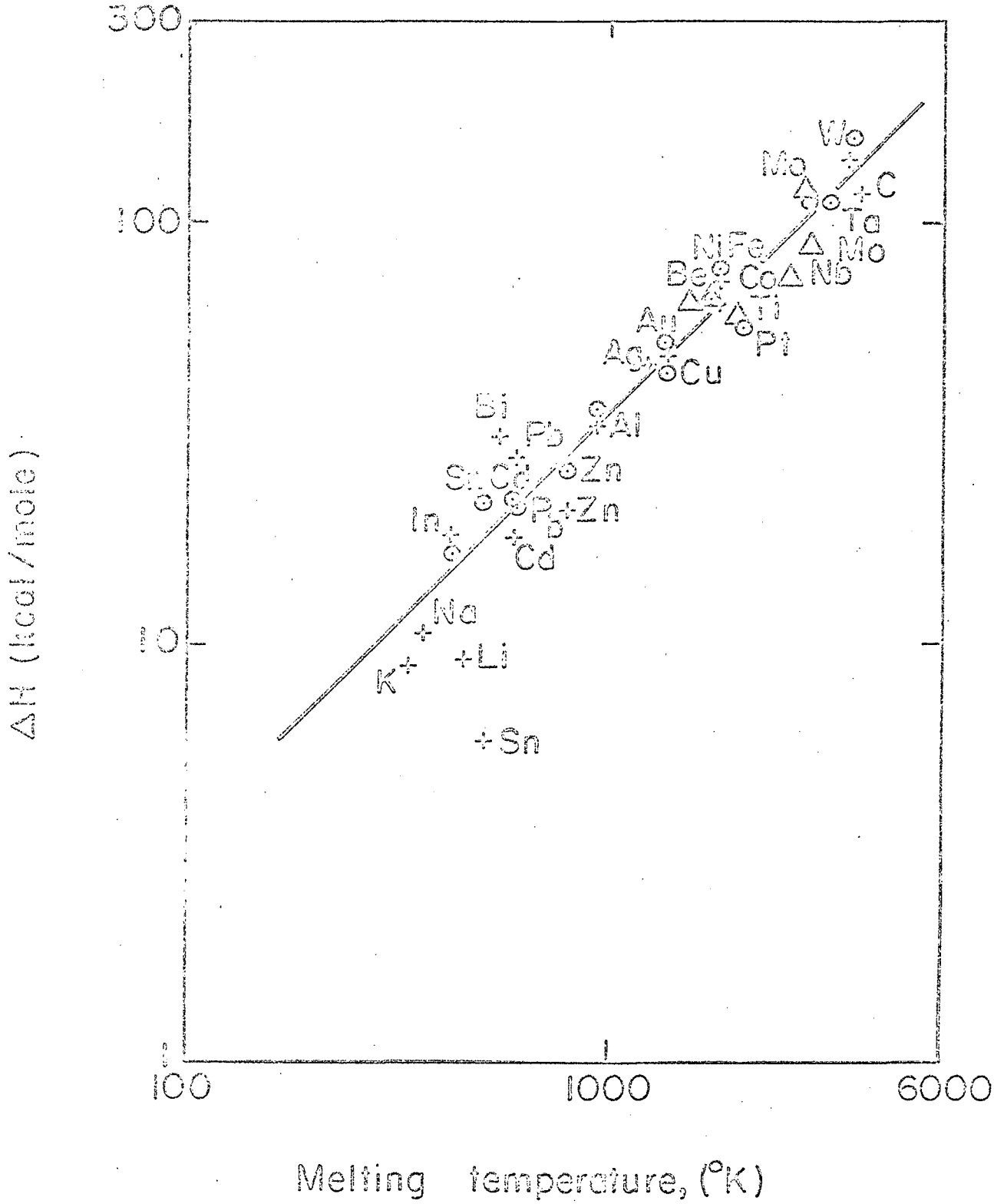
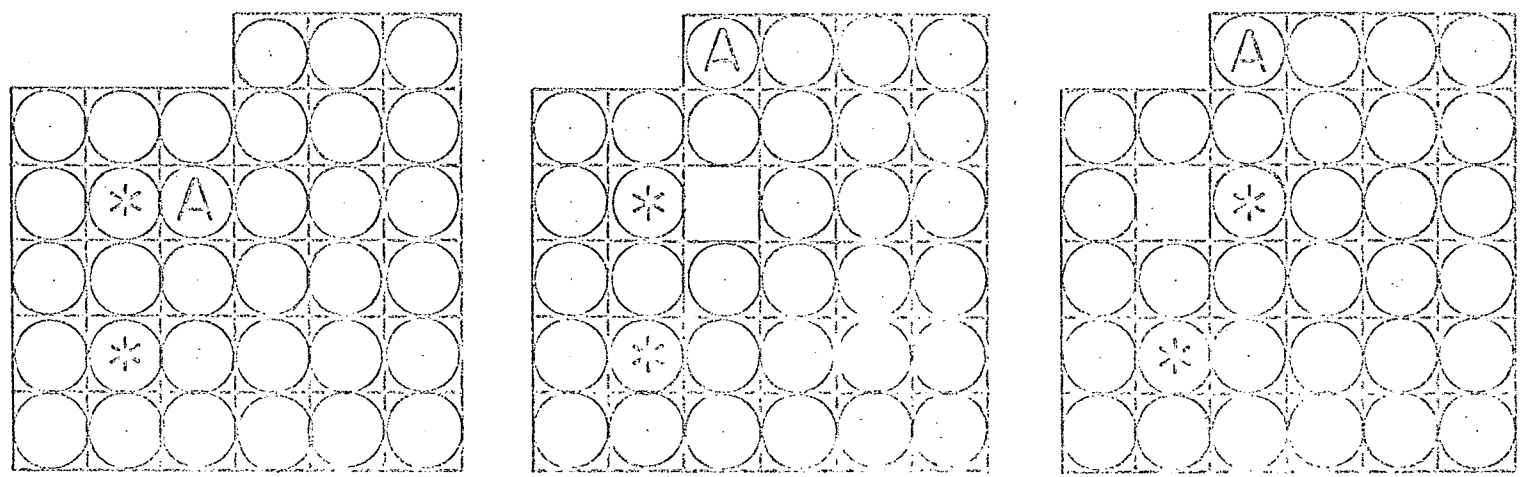


Fig. 7 Correlation of the Activation Energies for Creep, Stress-Rupture and Self-Diffusion with the Melting Temperature.[17]



(a) The production of a vacancy

(b) Radioactive * atom adjacent to a vacancy

(c) Atom vacancy exchange or motion of a vacancy

Fig. 8 Vacancies and Diffusion

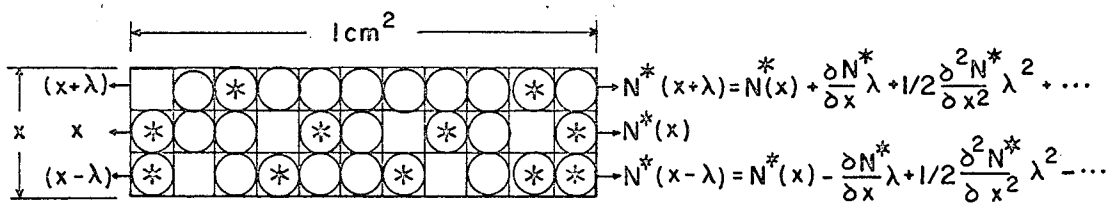


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Fig. 9 Diffusion of Tracer Atoms

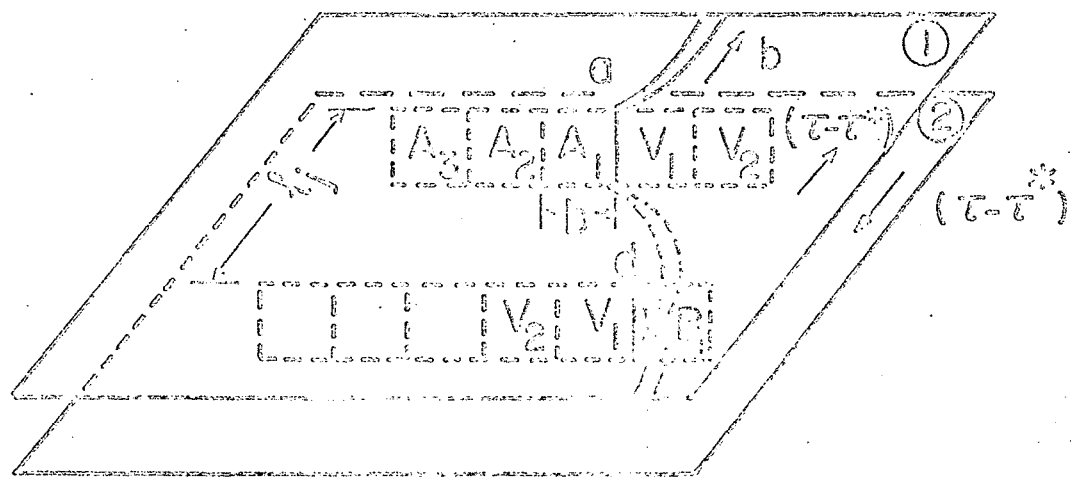


Fig. 10. Creation of a Vacancy by Motion of a Jog

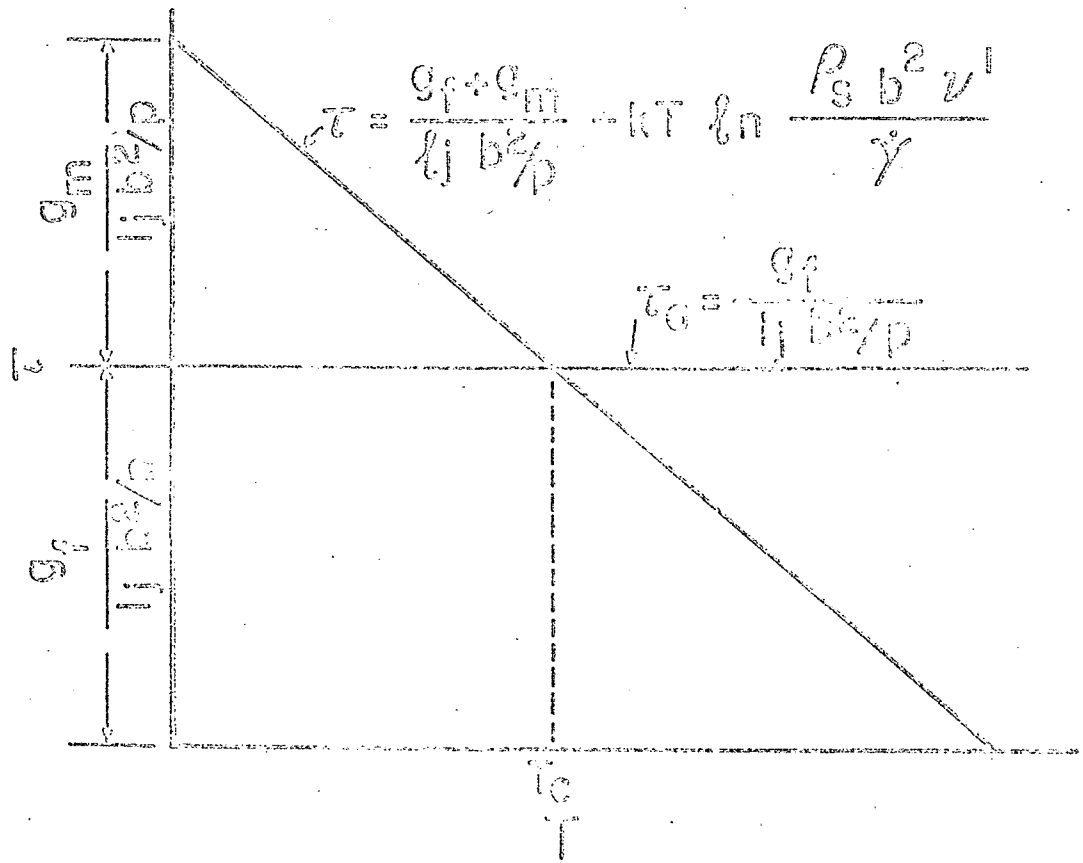


Fig. 11 Effect of Temperature of the Flow Stress for the Motion of Jogged Screw Dislocations

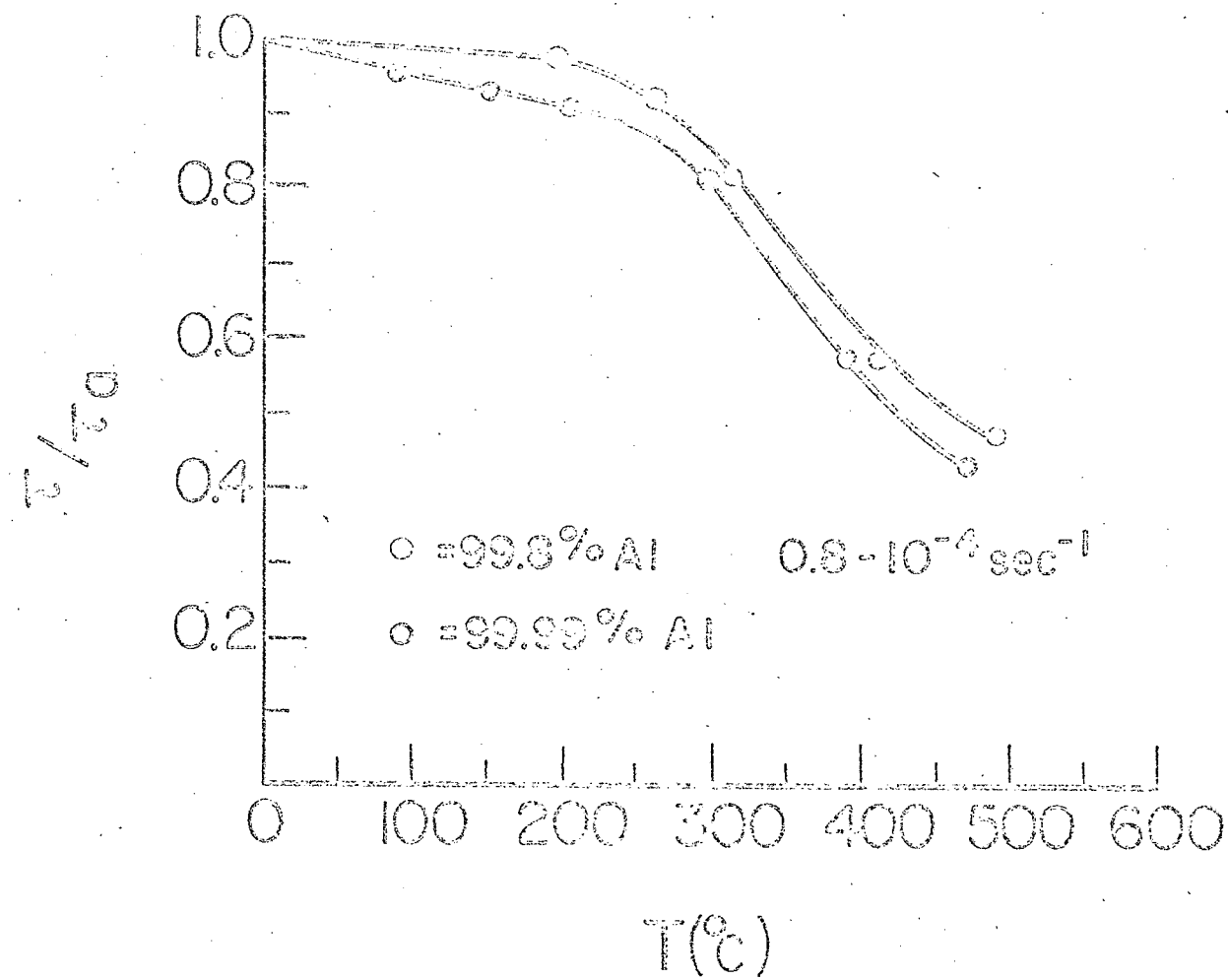
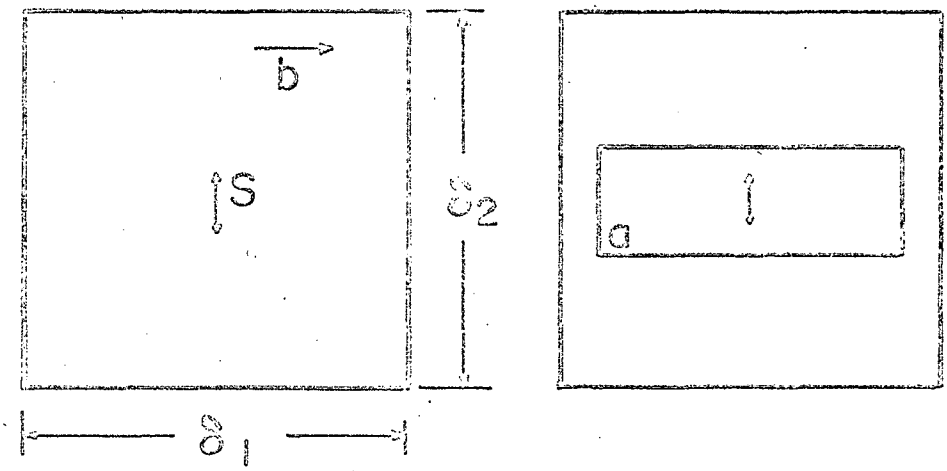
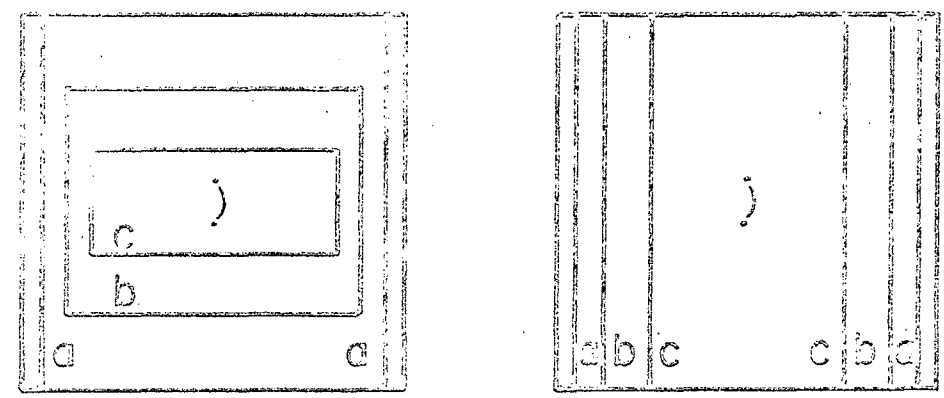


Fig. 12 Temperature Dependence of τ/τ_a for Two Purities of Polycrystalline Al (99.8%) and (99.9%)



(a) Dislocation source in center of subgrain to be formed upon deformation.

(b) Edge dislocations move rapidly to barriers which form subgrain boundary.



(c) Screw dislocations move slowly and cross slip.

(d) Edge dislocations pile up and stop source from operating, because of back stresses.

Fig. 13 Motion of Jogged Screw Dislocations

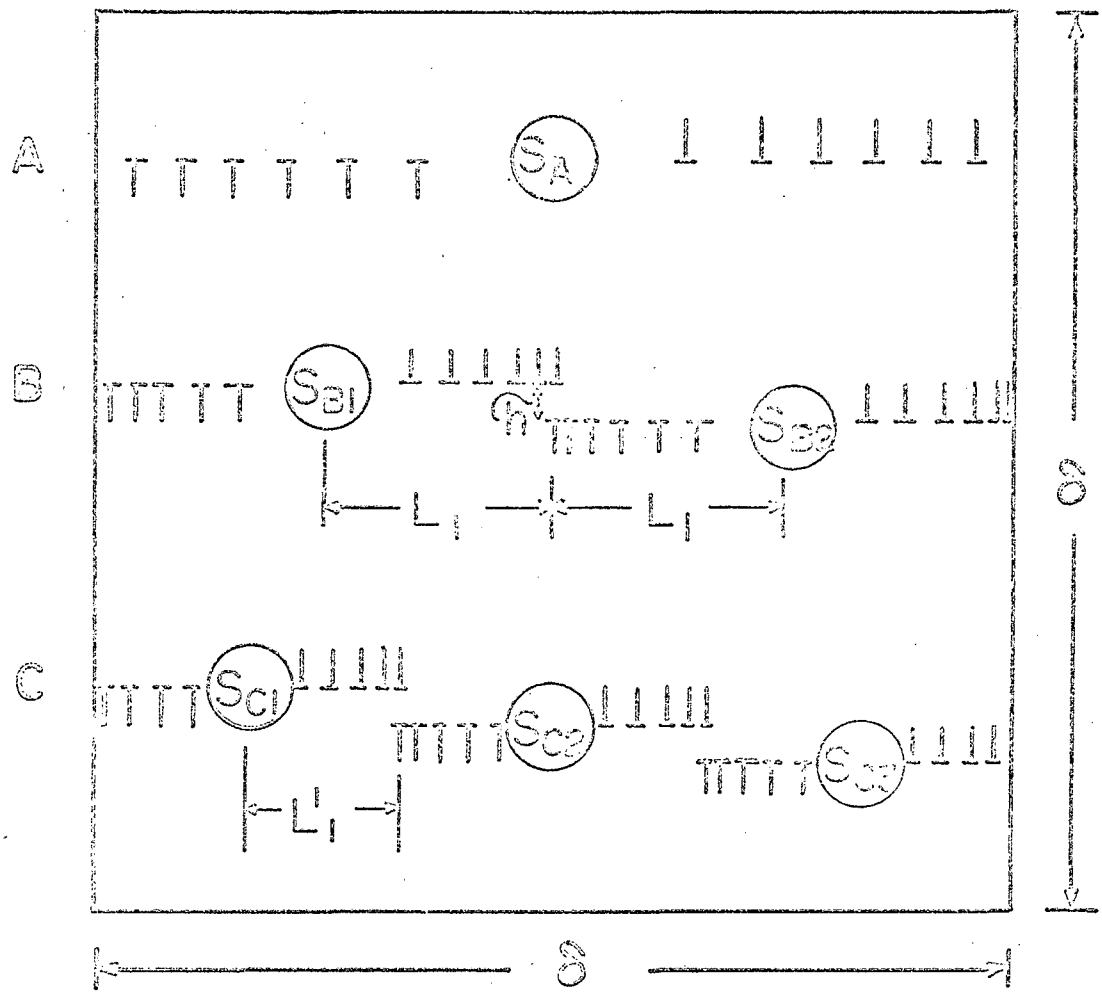


Fig. 14 Schematic Diagram of Possible Dislocation Arrays

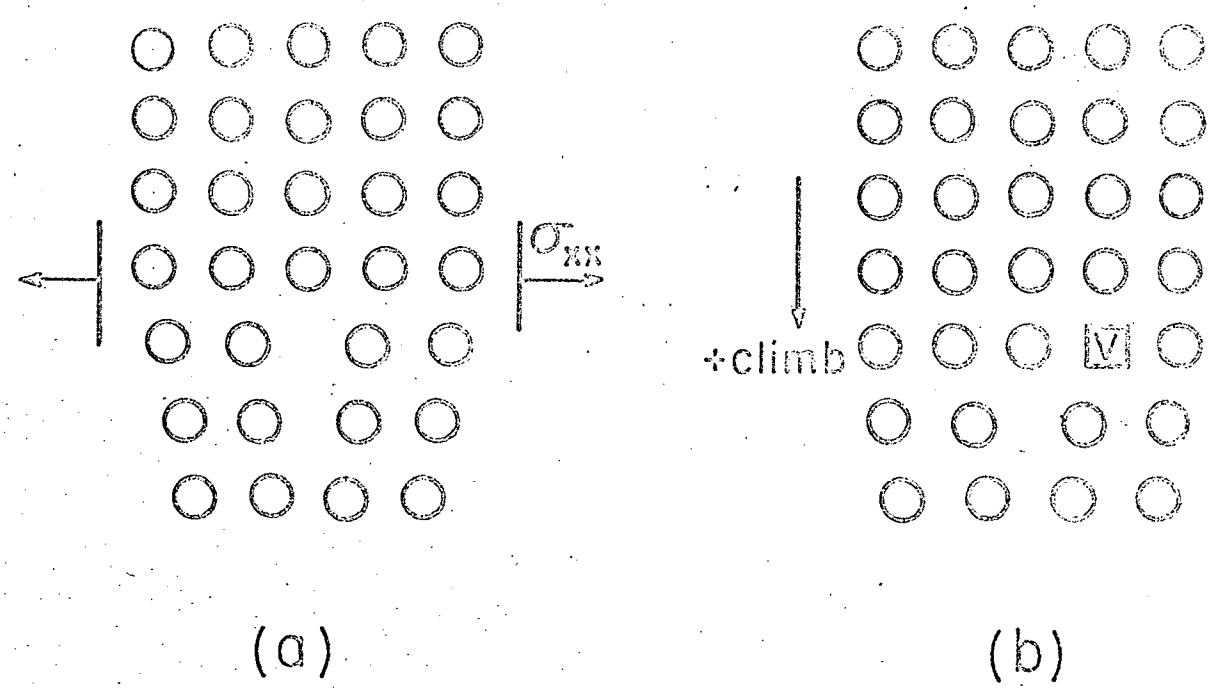


Fig. 15 Climb of an Edge Dislocation

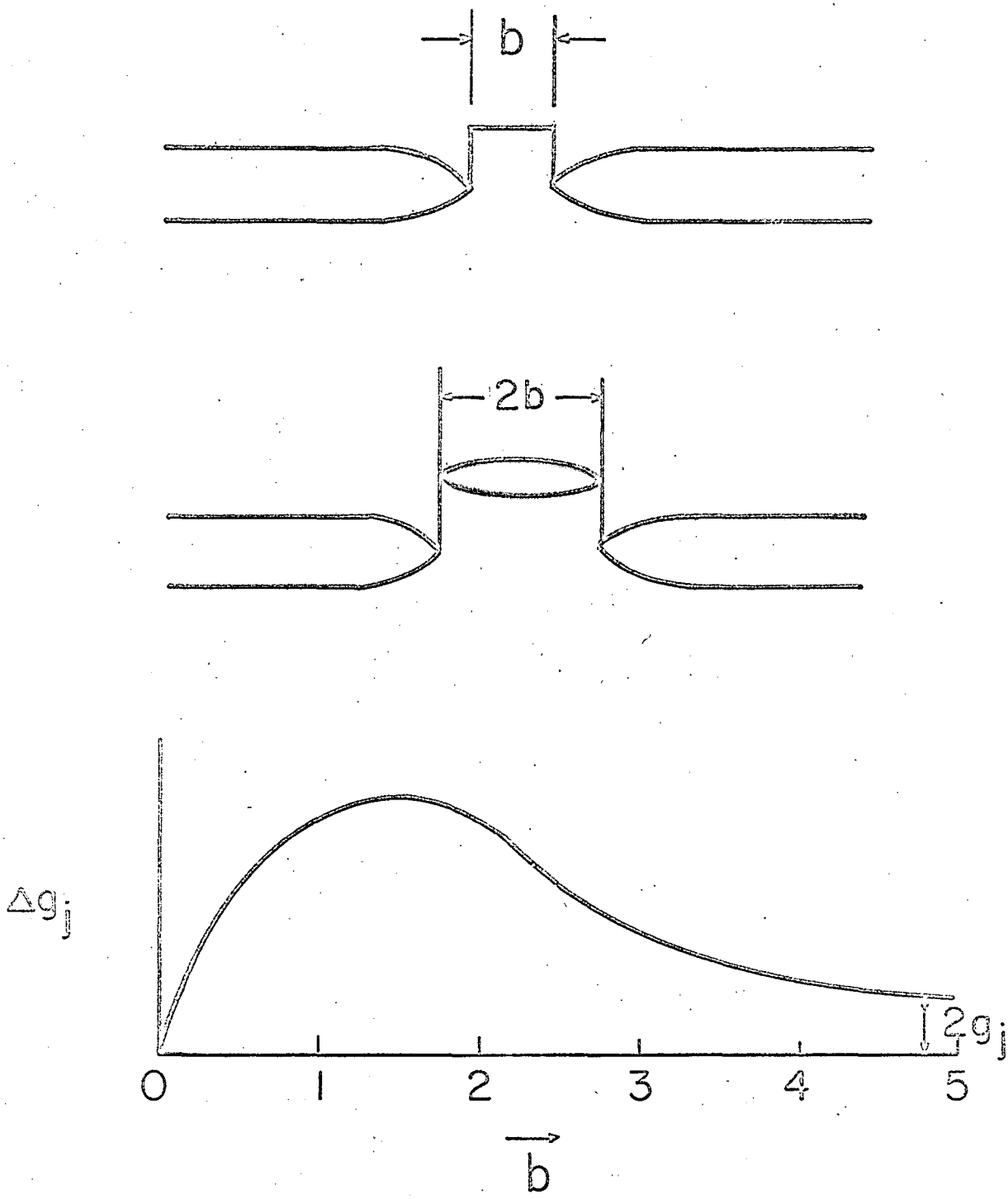


Fig. 16 Nucleation of Jogs

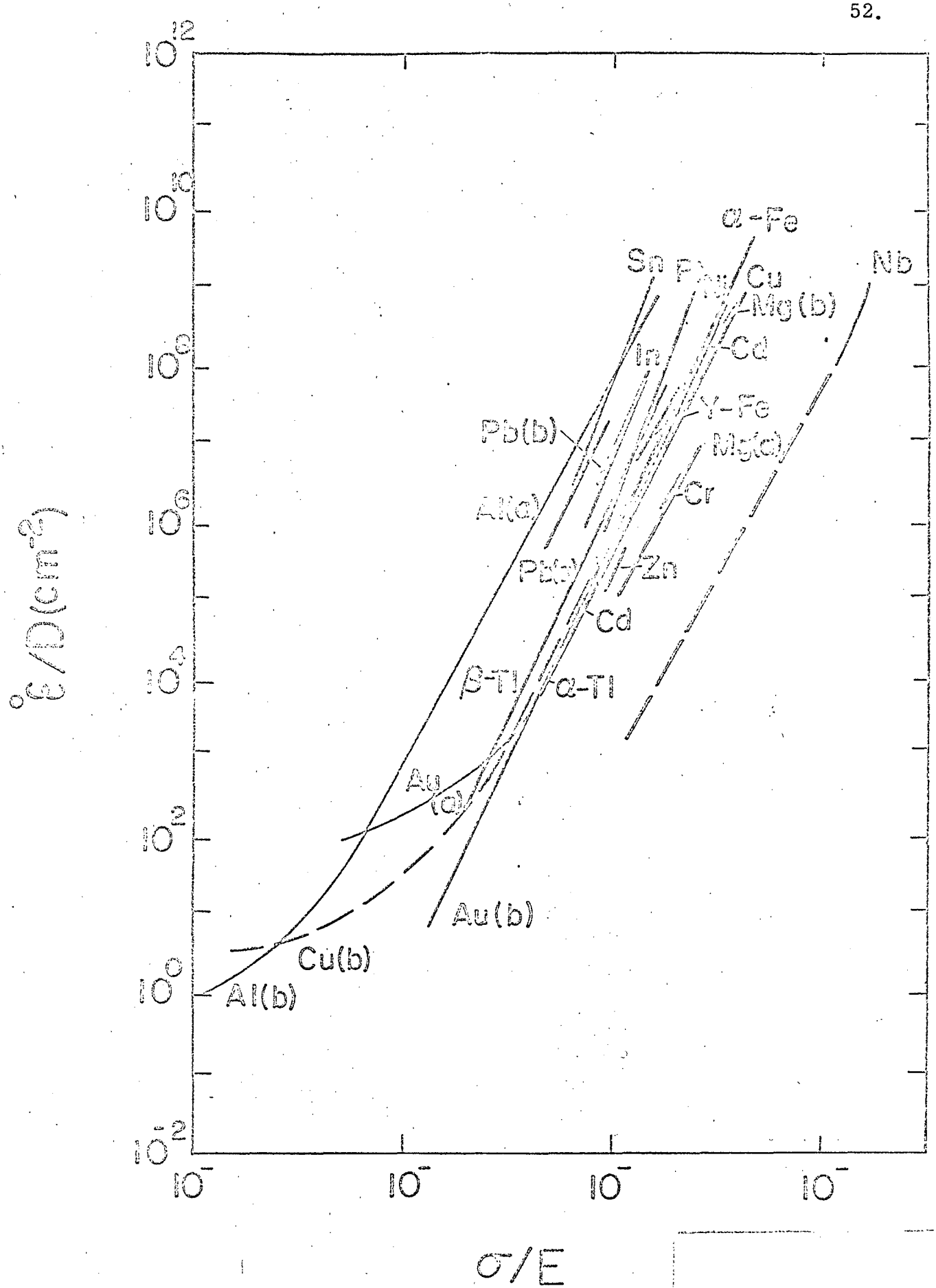


Fig. 17 Steady-State Creep Rate Provided by Diffusion Rate vs. Creep Stress Divided by Elastic Modulus for Various Pure Metals

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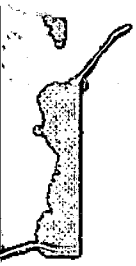
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