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### Publication Date

1965-12-01

UCRL-16562

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Lawrence Radiation Laboratory  
Berkeley, California  
AEC Contract No. W-7405-eng-48

MECHANISMS OF CREEP IN ORDERED ALLOYS

by

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December, 1965

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Written Specifically for the

"Folke Odquist Volume"

under

"Recent Progress in Applied Mechanics"

To Be Published in Honor of

Folke Odquist

at the Time of His Retirement as

Professor at the Royal Institute of Technology in Sweden

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## I. PREFACE

Professor Folke Odquist is known throughout the world for his outstanding contributions to applied mechanics and the behavior of materials. With great depth of perception and an imaginative mind he successfully pioneered research and development on all fronts of his chosen field. His vigor was matched only by his versatility. Creep of metals was one of the areas that he found intriguing. It is particularly fitting, therefore, that this paper, written in honor of his past works and in anticipation of his future contributions, describe some of the recent advances that have been made in rationalizing the creep behavior of ordered alloys.

## II. INTRODUCTION

High temperature creep in pure metals and random solid-solution alloys is now known to arise from the thermally-activated diffusion-controlled mechanisms of (a) stress-directed diffusion of vacancies (Nabarro, 1948), (b) the climb of arrested edge dislocations (Weertman, 1958), and (c) the non-conservative glide of jogged screw dislocations (Dorn and Mote, 1963). Whereas mechanism "a" is only significant near the melting temperature,  $T_m$ , the remaining diffusion-controlled mechanisms become facile at temperatures slightly in excess of about  $0.5 T_m$ ; above this point the creep resistance decreases rapidly with increasing temperature. Creep resistant engineering alloys are invariably used in a precipitation-hardened condition to provide additional high temperature strength by blocking the motion of dislocations. Unless the precipitated phase is

extremely insoluble, however, it will coarsen rather rapidly above  $0.5 T_m$  (Schoeck, 1961), attended with a corresponding decrease in creep resistance. On the other hand many ordered alloys retain high strengths up to from  $0.6$  to  $0.8 T_m$  (Westbrook, 1960). Furthermore, ordered alloys should not exhibit the high-temperature structural instability inherent in precipitation-hardened systems. In spite of the excellent potentialities of ordered alloys for creep-resistant application, relatively little has been done to foster their use. Perhaps the most significant area of research concerns the unraveling of the origin of their high-temperature strength. Once this is known it should be possible to suggest alloy compositions to achieve desired engineering properties. It is the purpose of this paper to describe some of the factors responsible for the high-temperature strength of ordered alloys.

### III. ORDER-DISORDER PHENOMENA

The unit cell of a simple type of ordered alloy is illustrated in Fig. 1. For complete order all of the A atoms reside on I sites and all of the B atoms on II sites. The more usual arrangement is characterized by disorder where each atomic species assumes random positions on the two types of sites. Ordered alloys are obtained only when the energy for the ordered state is less than that for the disordered state. The details of the various factors such as the Fermi energy, ionic interactions, strain energy, etc. that contribute to the difference in the energies of the ordered and disordered states have not yet been satisfactorily clarified. Modest progress in rationalizing order-disorder phenomena, however, has

been made from the somewhat naive chemical approach based on nearest neighbor atomic bonding. For the bond reaction  $A-B \rightarrow 1/2 A-A + 1/2 B-B$  the energy change is

$$\epsilon = \frac{\epsilon_{AA} + \epsilon_{BB}}{2} - \epsilon_{AB} \quad (1)$$

where  $\epsilon_{AB}$  is the energy of an A-B bond, etc. Only when  $\epsilon$  is positive can the ordered state be obtained.

Complete ordering is possible only when the atomic fractions  $N_A$  and  $N_B$  are equal to the fractions of the sites,  $N_I$  and  $N_{II}$  respectively. Nevertheless the tendency to ordering exists, although to a lesser degree, in compositions that deviate somewhat from the stoichiometric value. Frequently the degree of order,  $n$ , is defined as the fraction of atoms on the "right" sites minus the fraction of atoms on the "wrong" sites. Thus when all of the atoms are on "right" sites,  $n = 1$ , whereas when the atoms assume random positions  $n = 0$ . Bragg and Williams (1934) have shown that the energy per bond for the ordered alloy is  $\frac{-n^2\epsilon}{4}$  from that of the disordered state and the degree of order at thermal equilibrium is (for  $N_I = N_{II} = 1/2$ )

$$\frac{(N_A + n/2)(N_B + n/2)}{(N_A - n/2)(N_B - n/2)} = e^{\frac{-16n\epsilon}{kT}} \quad (2)$$

where  $kT$  is the Boltzmann constant times the temperature. For the equimolar alloy  $N_A = N_B = 1/2$  the degree of order vanishes at the Curie temperature,  $T_C = \frac{2\epsilon}{k}$ . The effect of temperature on the degree of order is shown in Fig. 2. Since the degree of order decreases very rapidly somewhat above  $T_C/2$ , only those alloys for which  $T_C > T_m$  will retain a

substantial degree of order in the creep range above  $T_m/2$ .

#### IV. DISLOCATIONS AND ANTIPHASE BOUNDARIES

Because the energy of a dislocation increases as the square of its Burgers' vector, the Burgers' vector is generally the shortest identity distance in the lattice. For random bcc alloys (Fig. 1) this is vectorially  $b = \frac{a}{2} [111]$  i.e. the distance from the center atom to any one of the corner atoms of the unit cell. For the completely ordered alloy, however, this is  $a[100]$  i.e. any edge of the unit cell. If, however, an  $a[111]$  dislocation dissociates as follows:  $a[111] = a[100] + a[010] + a[001]$  the energy change is zero; thus the  $a[111]$  superdislocation might exist. When, as described by Rachinger and Cottrell (1956), the bonding is substantially metallic, the  $a[111]$  dislocation will dissociate into an  $a/2[111] + a/2[111]$  with a decrease in energy proportional to  $-3/2a^2$ , thus leading to the most stable condition of a pair of  $a/2[111]$  dislocations. A schematic diagram of a dissociated superdislocation is shown in Fig. 3. As a result "wrong" bonds are formed causing an increase in energy. This energy is the antiphase boundary (APB) energy,  $\sigma$ , per unit length of the dislocation on the slip plane times the distance,  $y$ , of separation of the pair of dislocations. Thus there is a force  $\sigma$  per unit length of the dislocations causing them to coalesce. This is counterbalanced by the mutual repulsion between two similar dislocations, namely about  $\frac{Gb^2}{4\pi y}$ , where  $G$  is the shear modulus of elasticity (Dorn, 1963). Consequently dissociation is not complete and the equilibrium spacing between the two halves of the superdislocation in a common slip plane is about

$$y_0 = \frac{Gb^2}{4\pi\sigma} \quad (3)$$

Thus the two halves of the superdislocation are only slightly separated in alloys having high APB energies. In this event the two halves of a superdislocation will glide as a pair. When  $\sigma$  is not so high, the two halves of the superdislocation will be widely separated and each of the halves may glide more or less independently.

Even in the absence of dislocations, grains contain APBs. If one part of a grain nucleates to give the arrangement in Fig. 1 other parts may with equal probability nucleate so as to have the B atoms prefer the I sites and A atoms the II sites. Both domains will be equally highly ordered and each domain will be separated from the other by an antiphase boundary. In the simple CsCl type of alloy considered here only one type of phase boundary exists. Such simple boundaries must either terminate on the grain surface, or separate an enclosed domain from the surrounding domain. Since they have high energies, the antiphase boundary area per unit volume decreases when the alloy is annealed at elevated temperatures. This occurs rapidly in the CsCl type of alloy due to its simple structure and single type of domain boundary (Marcinkowski, 1963). In more complex alloys, however, the growth of domains is slow because a number of different kinds of domain boundaries can exist in metastable configurations as a foam structure (Marcinkowski, 1963).

The energy of an APB depends on the type of boundary and its orientation. Examples of the issues involved are illustrated by two dimensional models in Fig. 4. The APB in Fig. 4A was obtained by removing

the atoms along plane  $c$  ( $31\cdot$ ) and shifting the right half of the crystal by the antiphase vector  $1/2[\bar{1}\bar{1}\cdot]$ . Since equal numbers of A and B atoms were removed, no local compositional changes are encountered. Such APBs will be termed conservative. The A-B bonds that previously existed are now replaced by equal numbers of A-A and B-B bonds across this conservative APB. Two distinctly different APBs along the same ( $21\cdot$ ) plane are illustrated in Fig. 4B and 4C respectively. The plane  $c$  that was removed to form the APB contained either all A or all B atoms. Such APBs will be called non-conservative since the local composition across such APBs is no longer the average in the crystal. For the case shown in Fig. 4B where a plane of B atoms was removed, each four A-B bonds across the APB are replaced by three A-A bonds and one B-B bond; and in the example shown in Fig. 4C where a plane of A atoms was removed each four A-B bonds across the APB are replaced by one A-A bond and three B-B bonds.

Nonconservative APBs will not be readily produced. As two adjacent dissimilar domains nucleate, grow and impinge on each other the composition between them remains the average value. Consequently they would tend to form conservative APBs. When, due to local fluctuations in composition, a small nonconservative APB is produced another nonconservative APB must be produced nearby that compensates for the local compositional differences. Furthermore the free energy of a crystal will be greater when it contains nonconservative APBs. Consequently such nonconservative APBs will tend to form conservative APBs whenever facilitated by diffusion.

Conservative APBs are produced only on specific crystal planes. For example in the CsCl type of superlattice the A atoms occupy (000) sites and the B atoms ( $1/2, 1/2, 1/2$ ) sites. Crystal planes going through

(000) sites are given by  $hx + ky + lz = 1$ . If the plane is to contain B atoms as well, the general point  $p + 1/2, q + 1/2, r + 1/2$  where  $p, q, r$  are integers must also lie in the plane. Therefore

$$h(2p + 1) + k(2q + 1) + l(2r + 1) = 2 \quad (3)$$

Thus those planes for which  $h + k + l$  is an even number, contain equal numbers of A and B atoms and form conservative boundaries. Planes for which  $h + k + l$  is odd have two types of nonconservative boundaries, one through A atoms and the other through B atoms.

The total Burgers' vector for the CsCl lattice is  $a[111]$ . The possible slip planes must contain the Burgers' vector and are therefore given by the condition  $h + k + l = 0$ . Since the sum of indices is even all possible slip planes are planes that have conservative APBs. The climb, plane, for the  $a[111]$  Burgers' vector is the  $(111)$  plane. Dislocation climb consequently takes place on a nonconservative plane.

The unit surface energy,  $\sigma$ , of an APB is the excess energy of the bonds across the unit area of an APB above that which obtains in the crystal in the absence of an APB. To determine  $\sigma$  consider all bonds that cross the area  $a^2 \sqrt{h^2 + k^2 + l^2}$  which contains one lattice point. Let  $N$  be the total number of bonds from the sequence of planes below the APB that cross it. For conservative APBs, each A-B bond is replaced by  $1/2$  an A-A and  $1/2$  a B-B bond giving an energy change per bond of  $\epsilon$ . Therefore

$$\sigma = \frac{2\epsilon N n^2}{a^2 \sqrt{n^2 + k^2 + l^2}} \quad (4)$$

where the factor two reflects the fact that there are two atoms per lattice

site in the CsCl lattice and the average bond energy is acknowledged to depend on the square of the degree of order. Without loss of generality consider the plane for which  $h, k, \ell$  are all positive and  $h \geq k \geq \ell$ . The eight bond vectors issuing per atom in the CsCl type lattice are given by

$$t = a(\pm 1/2 \pm 1/2 \pm 1/2) \quad (5)$$

The number of planes  $n$  below an APB from which a bond can arise and cross the APB is equal to the projection of the bond vector in the direction normal to the plane of the APB divided by the distance between each plane. The unit vector normal to the plane is  $[hkl]/\sqrt{h^2 + k^2 + \ell^2}$  and the interplanar spacing is  $a/\sqrt{h^2 + k^2 + \ell^2}$ . Consequently

$$n = (\pm h/2 \pm k/2 \pm \ell/2) \quad (6)$$

The number of planes from which the various types of bonds can originate and cross a conservative APB are:

| Bond  | No. Planes          | condition         |
|-------|---------------------|-------------------|
| +++   | $(h + k + \ell)/2$  | all               |
| ++-   | $(h + k - \ell)/2$  | all               |
| + - + | $(h - k + \ell)/2$  | all               |
| + - - | $(h - k - \ell)/2$  | $h \geq k + \ell$ |
| - + + | $(-h + k + \ell)/2$ | $h \leq k + \ell$ |
| - + - | none                | all               |
| - - + | none                | all               |
| - - - | none                | all               |

Thus the total number of bonds issuing from all equivalent lattice sites

below the APB and crossing it is the sum of all above planes, whence  $N = 2h$  when  $h \geq k + l$  and  $N = (h + k + l)$  when  $h \leq k + l$ . Consequently, for conservative APBs in the CsCl type of lattice

$$\sigma = \frac{4h\epsilon n^2}{a^2 \sqrt{h^2 + k^2 + l^2}} \quad h \geq k + l \quad (7)$$

$$\sigma = \frac{2(h + k + l)\epsilon n^2}{a^2 \sqrt{h^2 + k^2 + l^2}} \quad h \leq k + l \quad (8)$$

The energy of a nonconservative APB can be determined by a modest extension of the previously given method for determining the energy of a conservative APB. In CsCl the nonconservative planes have alternately all A and all B atoms. The atomic interplanar spacing is one half of the lattice spacing and equals  $a/2\sqrt{h^2 + k^2 + l^2}$ . Thus the total number of planes  $n$  from which bonds can arise and cross an APB is given by

$$n = h + k + l \quad (9)$$

where now  $h + k + l$  is odd. Assume that the most distant plane from the APB that has bonds crossing the APB contains A atoms. Then the innermost plane must also contain A atoms. Consequently the number of planes from which bonds arise from A atoms is  $\frac{h + k + l + 1}{2}$  and the number from B atoms is  $\frac{h + k + l - 1}{2}$ . Therefore the numbers of planes from which various types of bonds arise and cross a nonconservative APB are as follows:

| Bond | No. Planes A Atoms  | No. Planes B Atoms  | Condition |
|------|---------------------|---------------------|-----------|
| +++  | $(h + k + l + 1)/2$ | $(h + k + l - 1)/2$ | all       |
| ++-  | $(h + k - l + 1)/2$ | $(h + k - l - 1)/2$ | all       |

|       |                      |                      |             |
|-------|----------------------|----------------------|-------------|
| + - + | $(h - k + l + 1)/2$  | $(h - k + l - 1)/2$  | all         |
| + - - | $(h - k - l + 1)/2$  | $(h - k - l - 1)/2$  | $h > k + l$ |
| - + + | $(-h + k + l + 1)/2$ | $(-h + k + l - 1)/2$ | $h < k + l$ |
| - + - | none                 | none                 | all         |
| - - + | none                 | none                 | all         |
| - - - | none                 | none                 | all         |

Thus the number of A-A and B-B bonds crossing the area per lattice point of a nonconservative APB are  $N_{A-A} = 2(h + 1)$  and  $N_{B-B} = 2(h - 1)$  when  $h > k + l$ , and  $N_{A-A} = (h + k + l + 2)$ ,  $N_{B-B} = (h + k + l - 2)$  when  $h < k + l$ . In the absence of an APB the number of A-B bonds crossing the same area is  $4h$  when  $h > k + l$  and  $2(h + k + l)$  when  $h < k + l$ . Thus, formally, the energy of this APB, obtained by removal of B atoms is

$$\sigma_B = \frac{N_{AA}\epsilon_{AA} + N_{BB}\epsilon_{BB} - N_{AB}\epsilon_{AB}}{a^2\sqrt{h^2 + k^2 + l^2}} \eta^2 \quad (10)$$

thus

$$\sigma_B = \frac{\{4h\epsilon + 2(\epsilon_{AA} - \epsilon_{BB})\}\eta^2}{a^2\sqrt{h^2 + k^2 + l^2}} \quad h > k + l \quad (11)$$

or

$$\sigma_B = \frac{\{2(h + k + l)\epsilon + 2(\epsilon_{AA} - \epsilon_{BB})\}\eta^2}{a^2\sqrt{h^2 + k^2 + l^2}} \quad h < k + l \quad (12)$$

If, however, the APB was made by the removal of a plane of A atoms the APB energy would have been

$$\sigma_A = \frac{\{4h\epsilon + 2(\epsilon_{BB} - \epsilon_{AA})\}\eta^2}{a^2\sqrt{h^2 + k^2 + l^2}} \quad h > k + l \quad (13)$$

or

$$\sigma_A = \frac{\{2(h + k + l) + 2(\epsilon_{BB} - \epsilon_{AA})\} \eta^2}{a^2 \sqrt{h^2 + k^2 + l^2}} \quad h < k + l \quad (14)$$

The above energies for nonconservative APBs must be applied with caution, because it is necessary to conserve types of atoms in the crystal. In general, therefore, nonconservative APBs come in pairs or are otherwise adjusted in number and extent so as to maintain the composition constant. When they appear in pairs the average surface energy of a nonconservative APB is given by the same equation as that for a conservative APB.

#### V. ENERGIES OF SUPERDISLOCATION PAIRS

At high temperatures each half superdislocation in edge orientation might climb from a slipplane so as to reside on a different parallel plane. Climb of the general type shown in Fig. 5A will be extremely difficult because of the increased energy required for the formation of jogs. Climb of the types depicted in Fig. 5B, however, is more readily possible especially when the two climb planes contain the same types of atoms and the local composition changes require diffusion over a short distance,  $r$ . Thus the net change in the area of nonconservative climb APB is zero. Each dislocation in Fig. 5B, however, now resides on different parallel slip planes. It is possible that diffusion can result in shifting the APB from plane XY to plane X'Y'. In Fig. 5C the dislocations might climb with relative ease onto a single new slip plane. Whether the climb as depicted in Figs. 5B and 5C will take place or not depends on the total energy of the various arrangements.

The energy of a superdislocation pair arises from their elastic interactions and the presence of the APB that joins them. Considering a unit long dislocation pair separated a distance  $r$  and joined by a planar APB, the energy due to the APB is  $\sigma r$  and the resulting force is  $-\sigma$  in the radial direction. Thus the sum of the elastic interaction forces and those arising from the APB, as given in polar coordinates are

$$F_r = \frac{A}{r} - \sigma \quad (15)$$

$$F_\theta = \frac{A}{r} \sin 2\theta \quad (16)$$

where  $A$  equals  $Gb^2/2\pi$  for a screw dislocation and  $Gb^2/2\pi(1 - \nu)$  for an edge dislocation,  $\nu$  being Poisson's ratio.  $\theta$  is always  $90^\circ$  for screw dislocations and is the angle between the Burgers' vector and the line lying in the APB that is perpendicular to the dislocation line. Thus  $\theta$  is zero for edge dislocations lying in their slip plane. Under these circumstances, the energy,  $U$ , of a superdislocation pair is

$$U = U\{b\} - \int_b^r A \frac{dr}{r} + \sigma r = U\{b\} - A \ln r/b + \frac{4hBr}{\sqrt{h^2 + k^2 + l^2}} \quad (17)$$

where  $U\{b\}$  is the elastic interaction energy when the two halves are a distance  $b$  apart,  $B = \epsilon\eta^2/a^2$  and  $h \geq k \geq l$ .

At equilibrium  $\partial U/\partial r = 0$  and

$$r = \frac{A}{B} \frac{\sqrt{h^2 + k^2 + l^2}}{4h} \quad (18)$$

Thus the energy for equilibrium spacing of the dislocations on a slip plane is

$$U = U\{b\} - A \ln \frac{A \sqrt{h^2 + k^2 + \ell^2}}{B 4hbe} \quad (19)$$

The effect of orientation on the energies of superdislocation pairs will be illustrated by taking the case of  $b = 1/2[\bar{1}11]$ . This dislocation exists on planes where  $-h + k + \ell = 0$  and the requirement that  $h \geq k + \ell$  is automatically satisfied when  $h$ ,  $k$ , and  $\ell$  are positive. The lowest energy is obtained for the (110) slip plane illustrating that superdislocations prefer to reside on the (110) slip planes in elastically isotropic ordered alloys of the CsCl type.

As shown by Flinn (1960), at high temperatures where diffusion is facile, edge dislocations might climb out of their slip plane. As shown by Eq. (16) the force on an edge dislocation in the  $\theta$  direction vanishes only for  $2\theta = 0$  (i.e. when both dislocations are on the same slip plane) and for  $2\theta = 180^\circ$ , (i.e. when the two edge dislocations are directly above each other and connected by a vertical APB. The pair of edge dislocations are in metastable equilibrium at  $\theta = 0$  since a small but finite force  $F_\theta$  will result for any small virtual change in  $\theta$  from zero. The pair of edge dislocations, however, will be in stable equilibrium when  $\theta = 90^\circ$ . Thus any displacement of one of a pair of edge superdislocations to a second slip plane, such as occurs from the production of jogs, would at high temperatures result in climb so that eventually the two halves of the edge dislocation would assume an arrangement of one above the other, the two being connected by a nonconservative APB.

## VI. BEHAVIOR OF ORDERED ALLOYS AT LOW TEMPERATURES

A brief account of some of the factors that affect the low temperature

plastic behavior of ordered alloys are presented here because they also play a role in high temperature creep. Whenever a superdislocation crosses an antiphase domain it produces a ledge on the boundary thereby increasing the energy (Cottrell, 1953). The work required to move a dislocation through an ordered alloy is therefore greater than that for a disordered alloy and the flow stress  $\tau_L$  to produce ledges is

$$\tau_L = \tau_G + \sigma/\ell \quad (20)$$

where  $\tau_G$  is the stress required to overcome interactions between dislocations, and  $\ell$  is the average distance a dislocation must travel to cut through a domain.

Antiphase boundaries are also produced as shown in Fig. 6 (Vidoz and Brown, 1962). Each half of glide superdislocation that has intersected a forest superdislocation in a screw orientation will acquire two unit jogs. The distance  $\ell_j$  between jogs will depend on the separation of the halves of forest superdislocations and on the geometric arrangements of the sets of dislocations upon intersection. The jogs on the trailing glide dislocation, however, will be displaced a distance  $d$  from those on the leading glide dislocation whenever the forest dislocation pair moves on their slip plane during the process of intersection. In this commonly expected event, two rectangular tubes of antiphase boundary will be produced. The distance  $\ell_T$  between pairs of tubes will depend on how many forest dislocations in screw orientation have been intersected.

When the glide dislocations are in screw orientation the jogs can slide along the dislocation conservatively; but since the jogs on each

half dislocation have the same sign they may, because of their mutual repulsion, move apart. Jogs on screw dislocations will, however, leave a trail of vacancies (or interstitials) behind them as the dislocation moves forward. Since the jogs on the trailing screw dislocation can slide, they will be pulled, due to the antiphase boundary energy, into alignment with the jogs on the leading dislocation during the forward motion of the dislocations thereby pinching off the tubes.

When the glide dislocation pair is in edge orientation, however, the jogs cannot slide conservatively along the dislocation. In these cases the tubes will not pinch off. Jogs on edge dislocations, however, can glide forward conservatively with the moving dislocation. In general a dislocation having both edge and screw components will exhibit an appropriate average behavior between that for pure screw and edge orientations. When the pairs of tubes are a distance  $l_T$  apart the stress required to move the dislocation is

$$\tau_T = \tau_G + \frac{4\sigma d}{2bl_T} \quad (21)$$

The total work done in moving a superdislocation forward will depend on the energy to produce both ledges and tubes. Therefore the flow stress is

$$\tau = \tau_G + \frac{\sigma}{l} + \frac{4\sigma d}{2bl_T} \quad (22)$$

Since  $l_T$  is very large at the initial stage of yielding, the yield stress is dictated primarily by the ledge mechanism. As deformation takes place, however,  $l_T$  decreases and the formation of tubes becomes the major

factor that determines the flow stress. Simultaneously, however, the alloy becomes more cluttered up with antiphase boundaries and the atomic arrangement approaches that of a random alloy causing  $\sigma$  to decrease with straining.

The statistical details of dislocation motion and multiplication are not yet well enough known to permit an accurate analysis of strain hardening in ordered alloys. The following admittedly crude formulation, however, is illustrative of the factors that are involved and the expected trends: Assuming that dislocations spread out from  $N_S$  sources per unit volume to a mean radius  $R$ , the shear strain is

$$\gamma = N_S \pi R^2 b \quad (23)$$

Then the distance between pairs of tubes is

$$l_T = \frac{2\pi R}{\alpha \pi R^2 \rho_f} = \frac{2}{\alpha R \rho_f} = \frac{2}{\alpha \rho_f} \sqrt{\frac{\pi N_S b}{\gamma}} \quad (24)$$

where  $\rho_f$  is the density of the forest dislocations and  $\alpha$  represents the fraction of those intersected that generate tubes. Under a fixed stress all dislocations will move until they acquire the same value of  $l_T$ . Under polyslip however,  $\rho_f$  will increase somewhat with strain but this effect is small and will be neglected here.

Assuming the domains are initially cubes of volume  $l_0^3$ , their initial surface area is  $6l_0^2$ . For each unit strain their surface area increases  $2l_0^2$  to give an area  $6l_0^2 + 2l_0^2\gamma$  as straining continues. Consequently the mean distance a dislocation must move to intersect the surface is about

$$\frac{1}{\ell} = \frac{\text{surface}}{\text{volume}} = \frac{(6\ell_o^2 + 2\gamma\ell_o^2)/6}{\ell_o^3} = \frac{1 + \gamma/3}{\ell_o} \quad (25)$$

Furthermore the density of pairs of tubes per unit volume also depends on the strain since

$$\rho_T = N_S \frac{2\pi R}{\ell_T} = N_S \alpha \pi R^2 \rho_f = \frac{\gamma \alpha \rho_f}{b} \quad (26)$$

The antiphase domain boundary area due to tubes and the domain surface is about

$$\rho_T \ell_o^3 4\ell_o d + \frac{1}{2} \left[ 6\ell_o^2 + 2\ell_o^2 \gamma \right] \quad (27)$$

where the factor 1/2 arises because only 1/2 of each boundary is ascribable to the domain in question. If every other slip plane were an antiphase domain the alloy would be substantially random. Thus the fraction of the alloy remaining ordered is

$$1 - \frac{2h}{\ell_o^3} \{ 4\rho_T \ell_o^4 d + 3\ell_o^2 + \ell_o^2 \gamma \} \quad (28)$$

where h is the interplanar spacing.

Therefore,

$$\sigma = \sigma_o \left\{ 1 - \frac{2h}{\ell_o} (4\rho_T \ell_o d + 3 + \gamma) \right\} \quad (29)$$

where  $\sigma_o$  is the antiphase boundary energy for the ordered alloy. When these factors are introduced Eq. (22) becomes

$$\tau = \tau_G + \sigma_o \left\{ 1 - \frac{2h}{\ell_o} \left[ \frac{4\gamma \alpha \rho_f \ell_o d}{b} + 3 + \gamma \right] \right\} \left\{ \frac{1 + \gamma/3}{\ell_o} + \frac{\alpha \rho_f d}{b} \left[ \frac{\gamma}{\pi N_S b} \right]^{1/2} \right\} \quad (30)$$

The yield strength ( $\gamma \approx 0$ ) increases as  $\ell_o$  decreases and exhibits a maximum value at  $\ell_o \approx 12h$ .

## VII. THERMALLY ACTIVATED CROSS SLIP

Cross slip from one to another slip plane can occur when the two slip planes have a common zone axis that is coincident with the orientation of a screw dislocation. Two alternate cross slip mechanisms can take place. If the two halves of a superdislocation are widely separated, as happens when the antiphase boundary energy is low, each might cross slip independently of the other. As each slips on the new glide plane it will leave an antiphase boundary in its wake. Consequently the stress to move such dislocations is athermal and given by  $\tau = \tau_G = \frac{\sigma}{b}$ . This is much higher than that required for the generation of tubes and would not be a common occurrence if the antiphase boundary energy were moderately high. When the two halves of the superdislocation, however, are only slightly separated, as occurs when the antiphase boundary energy is high, cross slip can occur in a manner analogous to cross slip for dissociated dislocations by the Friedel mechanism (1956) depicted in Fig. 7. The activation energy for cross slip is estimated to be about the constriction energy on the first plane plus that on the second plane (Stroh, 1954), namely

$$U_c = \frac{Gb^2 y_I}{20} \left( \ln \frac{y_I}{b} \right)^{1/2} + \frac{Gb^2 y_{II}}{20} \left( \ln \frac{y_{II}}{b} \right)^{1/2} \quad (31)$$

where  $y_I$  and  $y_{II}$  are the separations between the two halves of the superdislocation on planes I and II. Dislocations cross slip when obstacles arrest their forward motion on the original slip plane. Under these conditions, the stress  $\tau$  forces the trailing dislocation to approach the leading dislocation so that

$$\dot{\gamma}_I = \frac{Gb^2}{12\pi(\tau b + \sigma_I)} \quad (32)$$

where  $\sigma_I$  is the APB energy of plane I. In contrast

$$\dot{\gamma}_{II} = \frac{Gb^2}{12\pi\sigma_{II}} \quad (33)$$

since to effect cross slip there can be no obstacle in the vicinity of cross slip on plane II. Thus cross slip is facilitated by the stress and high values of  $\sigma_I$  and  $\sigma_{II}$ . To cross slip, segments of the two half superdislocations must first coalesce over a length of about  $6y$ . Consequently cross slip is unlikely to occur when the tubes are too close to each other (i.e.  $\ell_T < 6y$ ). Once a segment of a superdislocation has cross slipped, however, tubes can be dragged with the dislocation until their spacing reaches the limiting average value for the applied stress. The strain rate due to the cross slip mechanism is

$$\dot{\gamma}_{CS} = N \left( \frac{L}{6y} \right) A 2b \left( \frac{vb}{6y} \right) e^{-\frac{U_C}{kT}} \quad (34)$$

where  $N$  is the number of segments of length  $L$  held up at barriers,  $A$  is the area swept out on the cross-slip plane and  $v$  is the Debye frequency. It appears, however, that cross slip cannot be effective in facilitating creep so long as the extent of the motion of superdislocations on the cross slip plane is controlled by the tube forming mechanism.

#### VIII. HIGH TEMPERATURE CREEP

When a sufficiently high stress is applied to a crystalline material existing mobile dislocation and newly generated dislocations formed at points of stress concentration move forward and multiply. In ordered

alloys this forward motion is not only arrested by the usual interactions between dislocations as occurs in pure metals but also because the alloy becomes cluttered with APBs caused by formation of ledges at domain boundaries and as a result of formations of tubes due to the forward motion of jogged dislocations produced either by intersection or climb. At low temperatures the mechanical behavior is athermal. At moderately high temperatures, however, some thermally activated cross slip might cause some additional creep but, as previously described, this too will soon become negligibly small. Continued creep in ordered alloys is only possible at high temperatures as a result of diffusional processes that serve to eliminate the clutter of APBs. The various models for high temperature creep of ordered alloys are necessarily highly idealized simplifications of very complex situations. Consequently deductions based on such models do not always coincide exactly with the experimental facts. Nonetheless the models provide invaluable guides for rationalizing what factors are involved and their relative importance.

The creep rate  $\dot{\gamma}_i$  due to the  $i$ th process is given by

$$\dot{\gamma}_i = \rho_i b v_i \quad (35)$$

where  $\rho_i$  is the density of mobile dislocations undertaking this mechanism and  $v_i$  is the mean dislocation velocity (Dorn and Mote, 1963). If a series of mechanisms  $i, j, k, i, j, k$  etc. occur in sequence the creep rate is controlled by the slowest. Nevertheless, Eq. (34) yet applies to each process since each occurs to give the same rate; each so adjusts that  $\rho_i$

is the best for that process for which  $v_1$  is the greatest. When the processes are not sequential but are completely independent the net creep rate is the sum of those provided by each such independent process.

#### (A) Glide of Individual Dislocations

Consider first the creep rate of dissociated superdislocations each half of which has climbed or individually cross slipped to a different slip plane (vide Fig. 5B). For a length of one dislocation equal to  $b$  to move forward  $b$ , work equal to  $\sigma b^2$  must be supplied to make the APB. The stress, however, does work  $\tau b^3$  and therefore  $\sigma b^2 - \tau b^3$  must be supplied by a thermal fluctuation. But if this were all that took place the extra APB so produced would pull the dislocation back to its original position. To insure that the dislocation remains in its advanced position the extra APB produced must be destroyed by diffusion which occurs with about the Debye frequency  $\nu$  and requires the free energy  $g$  for diffusion. Thus the frequency of the forward motion of the dislocation is

$$v^+ = \nu Z e^{-\frac{(g + \sigma b^2 - \tau b^3)}{kT}} \quad (36)$$

where  $Z$  is the coordination number. Since the slip plane is conservative it contains equal numbers of A and B atoms and  $g$  can be approximated by the average of that for the two atomic species. The frequency for the reversed reaction, obtained similarly, is

$$v^- = \nu Z e^{-\frac{(g + \sigma b^2 + \tau b^3)}{kT}} \quad (37)$$

whence the mean velocity is

$$v_a = (v^+ - v^-)b = 2vbz e^{-\frac{(g + \sigma b^2)}{kT}} \sinh \frac{\tau b^3}{kT} \quad (38)$$

But  $\tau b^3/kT$  is usually small relative to unity and consequently

$$\dot{\gamma}_a = 2\rho_a b^2 v_z e^{-\frac{(g + \sigma b^2)}{kT}} \frac{\tau b^3}{kT} \quad (39)$$

Several factors deserve special mention: The gliding dislocations will usually be jogged. Under these circumstances the  $\sigma b^2$  term should be replaced by  $\sigma b^2 (1 + p_j)$  where  $p_j$  is the probability of finding a jog. Furthermore  $\rho_a$  will depend on the rates of cross slip and climb. Thus  $\rho_a$  will increase as the stress increases and it may also depend on the temperature.

#### (B) Glide of Superdislocation Pairs

As previously described, superdislocations moving as a pair on their slip plane become jogged and therefore form tubes of APBs as they move. The APB energy that must be supplied to move a pair of tubes forward one Burgers' vector is about  $\sigma b d$ . The stress, however, does work  $\tau \ell_T b^2$  during this forward motion. Thus the energy that must be supplied by a thermal fluctuation to produce one atomic area of an APB in the tubes is about

$$(4\sigma b d - 2\tau \ell_T b^2) b^2 / 4db = \sigma b^2 - \frac{\tau \ell_T b^3}{2d} \quad (40)$$

To destroy the APB formed in this manner requires diffusion having the free energy  $g$ . Thus the net frequency of the forward motion of the dislocation one Burgers' vector is about

$$v^+ = v \left( \frac{b}{\ell_T} \right) \left( \frac{b^2 Z}{4db} \right)^c e^{-\frac{(g + \sigma b^2)}{kT}} e^{-\frac{\tau \ell_T b^3}{2dkT}} \quad (41)$$

where the first term in the brackets corrects the Debye frequency for the frequency of vibration of the dislocation of length  $\ell_T$  and the second term corrects the Debye frequency so as to account for all atoms in a Burgers' vector length of the tubes. A reversed reaction of the same type occurs against the action of stress  $\tau$ . Consequently the net velocity of motion of the dislocation is

$$v = \frac{vb^3 Z}{2\ell_T d} e^{-\frac{(g + \sigma b^2)}{kT}} \sinh \frac{\tau \ell_T b^3}{2dkT} \quad (42)$$

and the creep rate is given by

$$\dot{\gamma}_p = \frac{\rho_p vb^4 Z}{2\ell_T d} e^{-\frac{(g + \sigma b^2)}{kT}} \sinh \frac{\tau \ell_T b^3}{2dkT} \quad (43)$$

### (C) Viscous Creep

Weertman (1957) has shown that, where dislocations from fixed sources have a velocity that depends linearly on the stress and arrested dislocations can be removed by facile climb,

$$\dot{\gamma}_C = \frac{2\pi(1 - \nu)Z}{G^2} \tau^3 B \quad (44)$$

where  $B$  is the mobility defined by  $v = B\tau b$ . Under the limiting conditions where either  $\tau b^3/kT \ll 1$  or  $\tau \ell_T b^3/2dkT \ll 1$  these conditions may be satisfied for (a) the glide of individual dislocations or (b) the glide of superdislocation pairs where

$$B_a = \frac{2vb^3Z}{kT} e^{-\frac{(g + \sigma b^2)}{kT}} \quad (45)$$

for glide of individual dislocations and

$$B_b = \frac{vb^5Z}{4d^2kT} e^{-\frac{(g + \sigma b^2)}{kT}} \quad (46)$$

for the viscous glide of dislocation pairs.

#### (D) Climb of Dislocations

In the CsCl lattice, climb takes place on the (111) planes which are nonconservative. The superdislocation  $a[111]$  is  $a\sqrt{3}$  long and the (111) lattice planes are  $a/\sqrt{3}$  apart. Consequently the superdislocation cuts across three lattice planes or six atomic planes of alternate layers of A and B atoms. A dissociated superdislocation in the CsCl lattice is shown in Fig. 8. When the APB energy is high the two halves of the superdislocation will be only slightly separated and therefore subjected to about the same local stress  $\sigma_{xx}$ . Under these conditions both halves of the superdislocation will climb together, and as climb takes place nonconservative APBs are produced on the (111) climb plane. The rate of climb will be about

$$v_c = vb p_v p_j Z e^{-\frac{(g_m - \sigma_{xx} b^3)}{kT}} - vb p_j Z e^{-\frac{(g_f + g_m + \sigma_{xx} b^3)}{kT}} \quad (47)$$

where  $p_v = e^{-g_f/kT}$  is the probability of a vacancy being present,  $p_j$  the probability of a jog being present,  $g_m$  the free energy of motion of a vacancy,  $Z$  is the coordination number and  $\sigma_{xx}$  is the stress normal to

the climb plane. The antiphase boundary energy does not enter this formulation because the two halves of the superdislocation move in unison and no extra APB need be produced. Thus the velocity of climb is

$$v_c = 2vb_p j Z e^{-g/kT} \sinh \frac{\sigma_{xx} b^3}{kT} \quad (48)$$

Under conditions of slow creep the superdislocations issuing from sources S are expected to take up positions as shown in Fig. 9, one of the pair of superdislocations being placed vertically above the other and connected with an APB. When the two arrays are close together the motion of the dislocations in the arrays is arrested. Further creep can only occur when the leading superdislocations of both arrays climb and annihilate each other. Upon annihilation another pair of superdislocations issue from the source. The creep rate due to this process for low stresses is crudely estimated to be (Dorn, 1963)

$$\dot{\gamma}_c = \frac{15,600 N L^5 Z \pi^4 \tau^5 b_p j}{(G/1-u)^4 kT} e^{-\frac{g}{kT}} \quad (49)$$

where N is the number of sources per cm<sup>3</sup> and L is defined in Fig. 9.

The quantity NL<sup>5</sup>, however, can also depend on the stress and therefore the creep rate need not coincide exactly with the fifth power law.

#### SUMMARY

The unique behavior of ordered alloys centers around the important effect of APBs on the motion of dislocations. Under high temperatures superdislocation pairs will climb so as to be vertically arrayed on

parallel slip planes and joined by APBs. Creep will then occur by glide of superdislocation pairs (vide VIII B). The activation energy for creep in ordered alloys exceeds that for self diffusion by  $ob^2$ , the energy to produce an atomic area of antiphase boundary. Furthermore the activation energy for diffusion will be slightly higher in the ordered than in the disordered condition, a factor which further enhances the creep resistance. Under steady state conditions, of course, creep is also determined by the climb process. But climb of superdislocations can be relatively rapid in contrast to glide. Consequently the steady state creep rate of ordered alloys is expected to follow Weertman's viscous creep behavior as has been observed in several cases (Lawley et al., 1960; Soler-Gomez and Tegart, 1964; Mukherjee and Dorn, 1964).

#### ACKNOWLEDGMENT

This report was prepared as part of the activities of the Inorganic Materials Research Division of the Lawrence Radiation Laboratory of the University of California, Berkeley, and was done under the auspices of the U. S. Atomic Energy Commission.

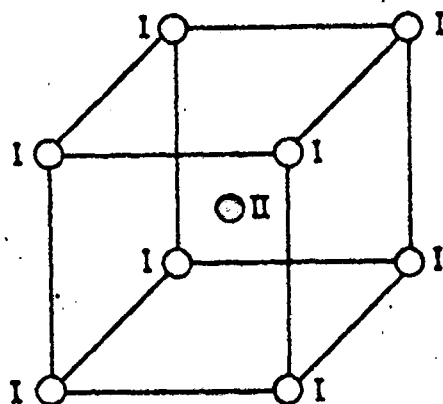
#### REFERENCES

1. Bragg, W. L. and Williams, E. J., Proceedings of the Royal Society, A145, 699 (1934).
2. Cottrell, A. H., Relation of Properties of Microstructure, A.S.M. Seminar, 131 (1953).
3. Dorn, J. E., "Energetics of Dislocation Mechanics", Proceedings of the National Research Council Seminar on Energetics in Metallurgical Phenomena (1963).
4. Dorn, J. E. and Mote, J. D., "Physical Aspects of Creep", High Temperature Structures and Materials, Proceedings of the Third Symposium on Naval Structural Mechanics, p.95-168, Pergamon Press, New York (1963).
5. Flinn, P. A. "Theory of Deformation of Superlattices, Trans. AIME, 218, 145-154 (1960).
6. Friedel, J., Phil. Mag., 46, 514, 1169 (1956).
7. Lawley, A., Coll, J. A. and Cahn, R. W. "Influence of Crystallographic Order on Creep of Fe-Al Solid Solutions", Trans. AIME, 218, 166 (1960).
8. Marcinkowski, M. J. "Theory and Direct Observation of Antiphase Boundaries in Superlattices", Electron Microscopy and Strength of Crystals, Interscience (1963).
9. Mukherjee, A. K. and Dorn, J. E., "On the Mechanisms of Creep in the Intermetallic Compound AgMg at Elevated Temperatures", J. Inst. of Metals, 93, 397 (1964).
10. Nabarro, F. R. N. "Deformation of Crystals by Motion of Single Ions", Report of a Conference on the Strength of Solids, London: Phys. Soc., 75-90 (1948).

11. Rachinger, W. A. and Cottrell, A. H., "Slip in Single Crystals of the CsCl Type", Acta Met., 4, 109 (1956).
12. Schoeck, G., "Thermodynamic Principles in High-Temperature Materials", Mechanical Behavior of Materials at Elevated Temperatures, 70-74, McGraw-Hill, New York, J. E. Dorn editor, (1961).
13. Soler-Gomez, A. J. R., and Tegart, W. J. "The High Temperature Creep of Ordered and Disordered MgCd", Acta Met., 12, 961 (1964).
14. Stroh, A. N., "Constrictions and Jogs in Extended Dislocations", Proc. of the Phys. Soc. London, A223, 404 (1954).
15. Vidoz, A. E. and Brown, L. M., "On Work-Hardening in Ordered Alloys", Phil. Mag., 1, 1167 (1962).
16. Weertman, J., J. Appl. Phys., 28, 1185 (1957).
17. Weertman, J., "Theory of Creep of Dispersioned-Hardened Alloys", Naval Research Laboratory Report No. 5123 (1958).
18. Westbrook, J. H., "Mechanical Properties of Intermetallic Compounds", 1-71, John Wiley and Sons Inc., New York (1960).

FIGURE CAPTIONS

1. Unit cell of CsCl type of ordered alloys.
2. Effect of temperature on the degree of long-range order.
3. Schematic diagram of a dissociated superdislocation.
4. Antiphase boundaries in the CsCl ordered lattice.
5. Possible climb configurations for superdislocations.
6. Formation of antiphase boundary tubes by motion of jogged superdislocations.
7. Cross slip of superdislocation segment.
8. Dissociated superdislocation  $a[111] \rightarrow \frac{a}{2} [\bar{1}\bar{1}1] + \frac{a}{2} [\bar{1}\bar{1}1]$ .
9. Equilibrium configuration of superdislocation arrays at high temperatures.



○ A-ATOMS  
○ B-ATOM

FIG. 1. UNIT CELL OF CsCl TYPE  
OF ORDERED ALLOYS.

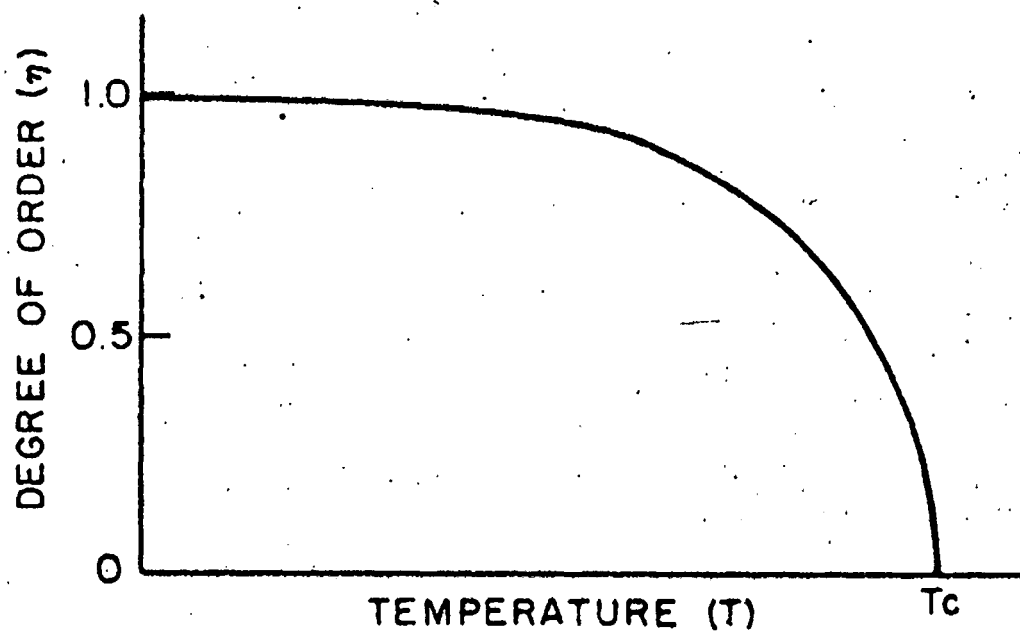


FIG. 2. EFFECT OF TEMPERATURE ON THE DEGREE OF LONG-RANGE ORDER.

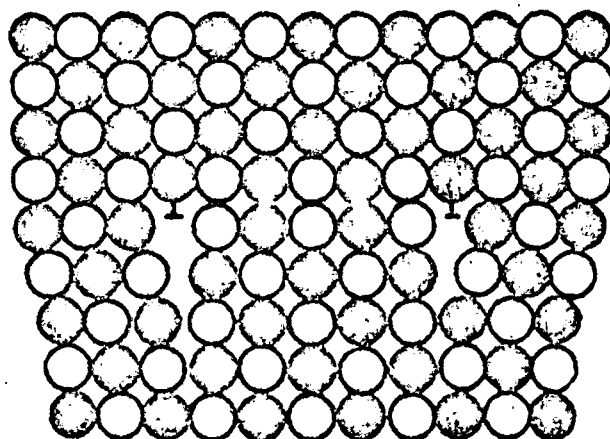


FIG. 3. SCHEMATIC DIAGRAM OF A  
DISSOCIATED SUPERDISLOCATION.

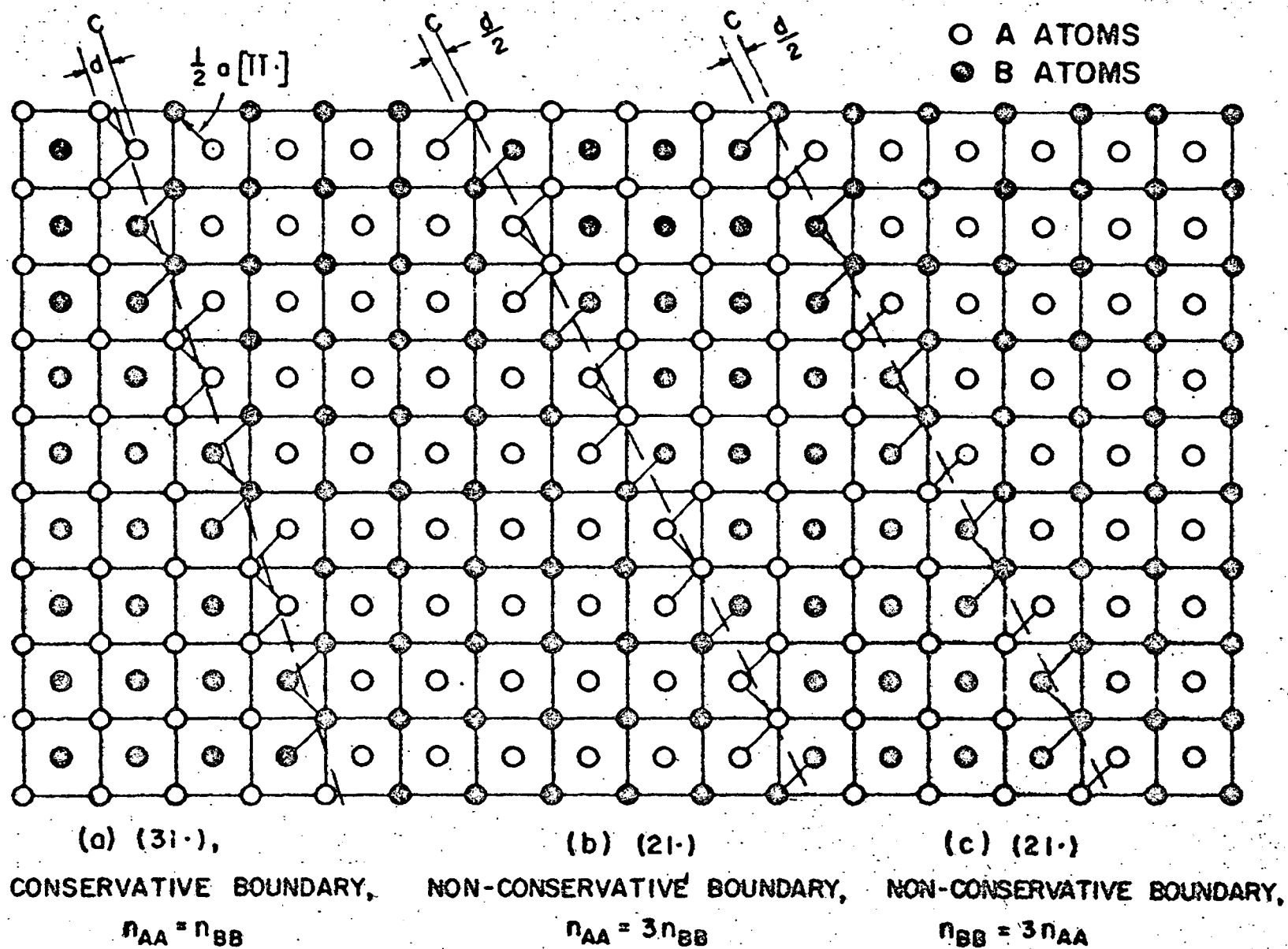
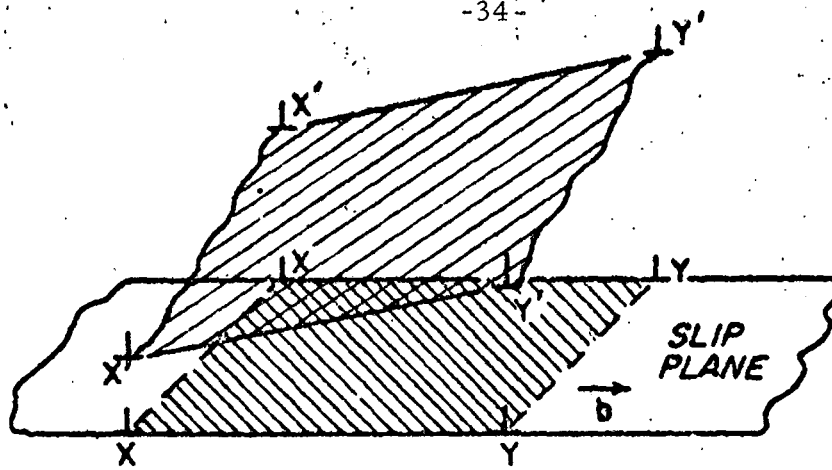
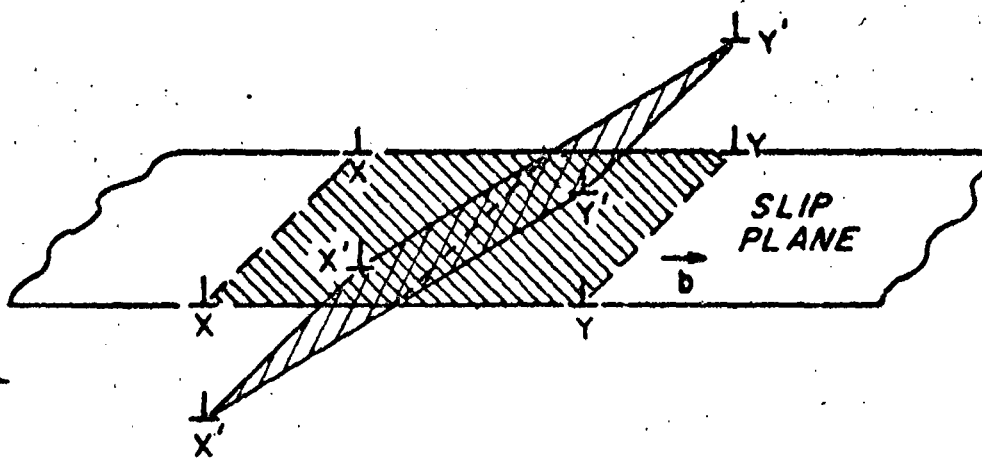


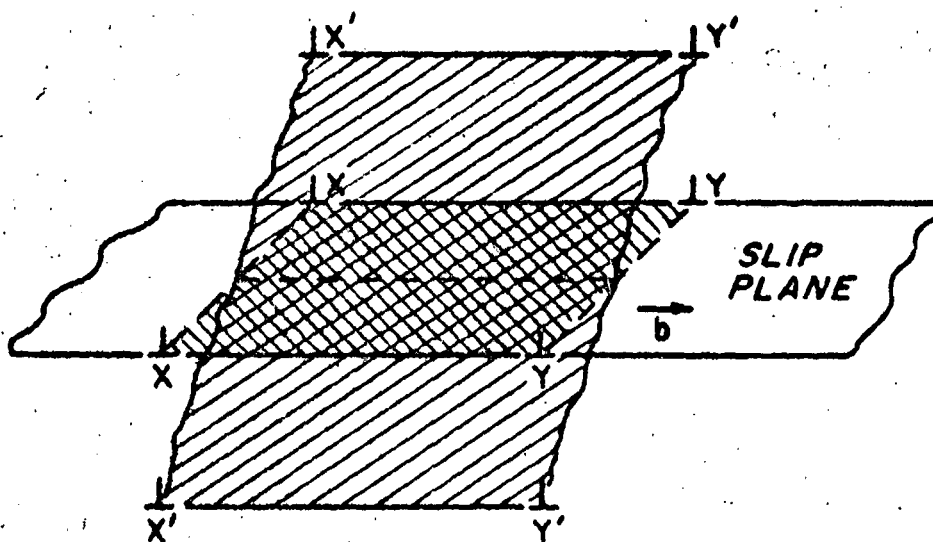
FIG. 4. ANTIPHASE BOUNDARIES IN THE CsCl ORDERED LATTICE.



(a)



(b)



(c)

FIG. 5 POSSIBLE CLIMB CONFIGURATIONS OF SUPERDISLOCATIONS.

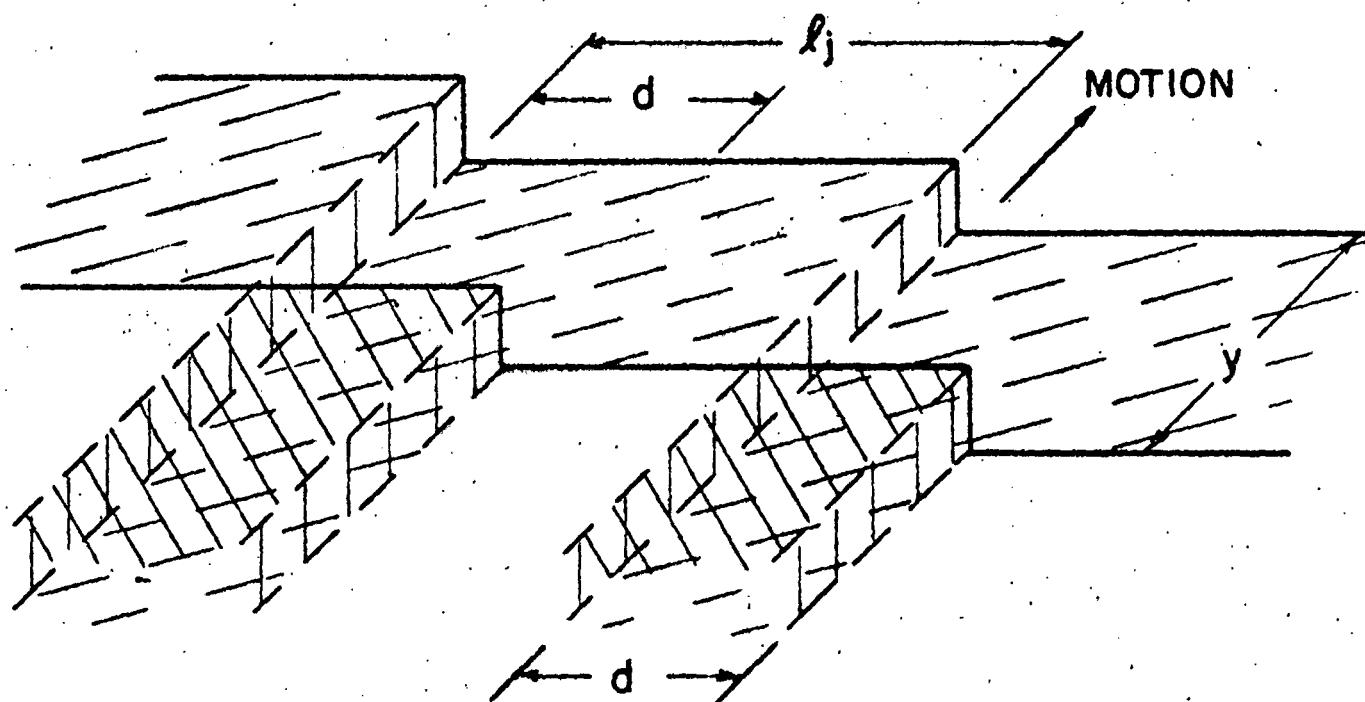


FIG. 6. FORMATION OF ANTIPHASE BOUNDARY TUBES BY MOTION OF JOGGED SUPERDISLOCATIONS.

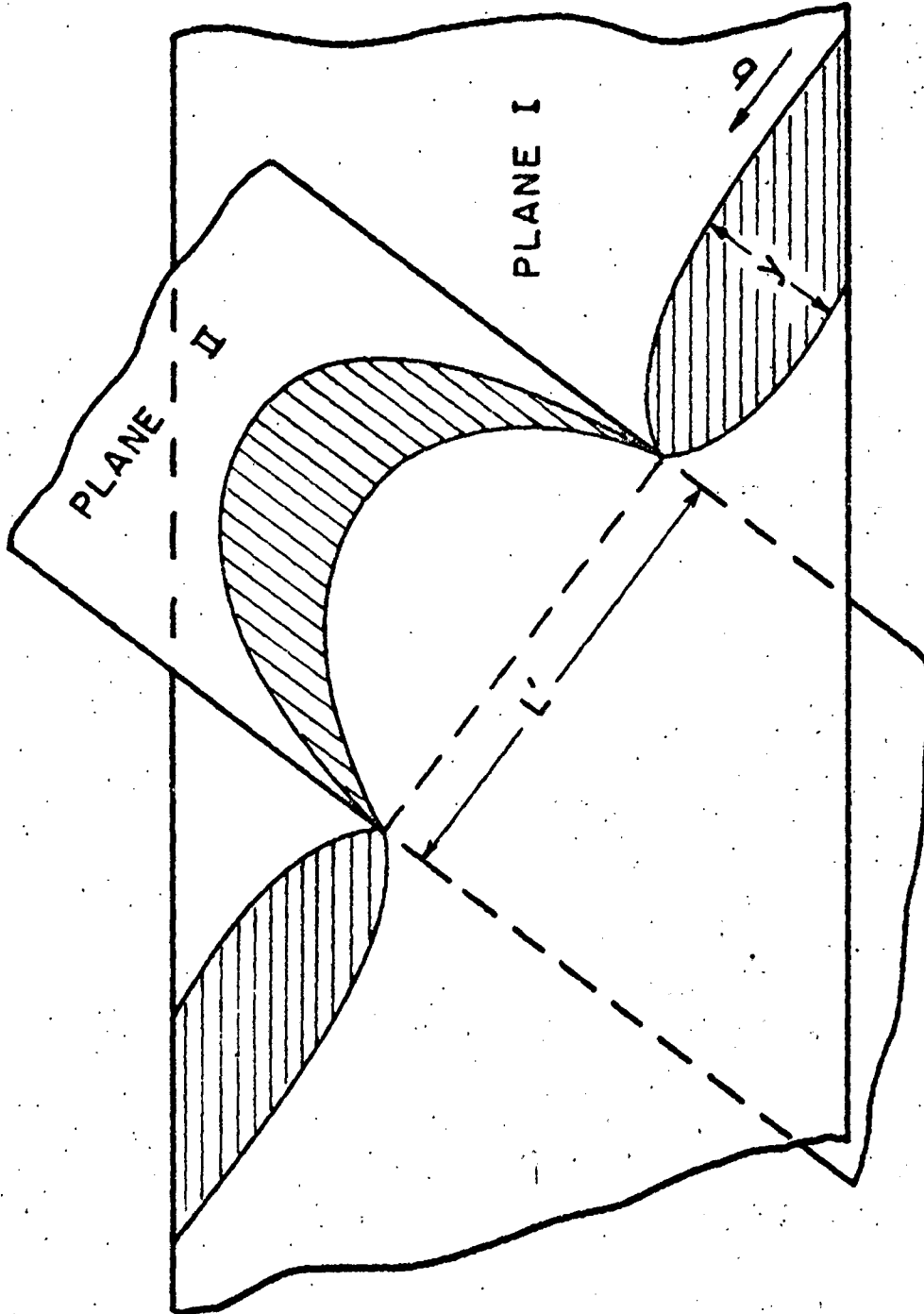


FIG. 7. CROSS SLIP OF SUPERDISLOCATION SEGMENT  
FROM PLANE I TO PLANE II.

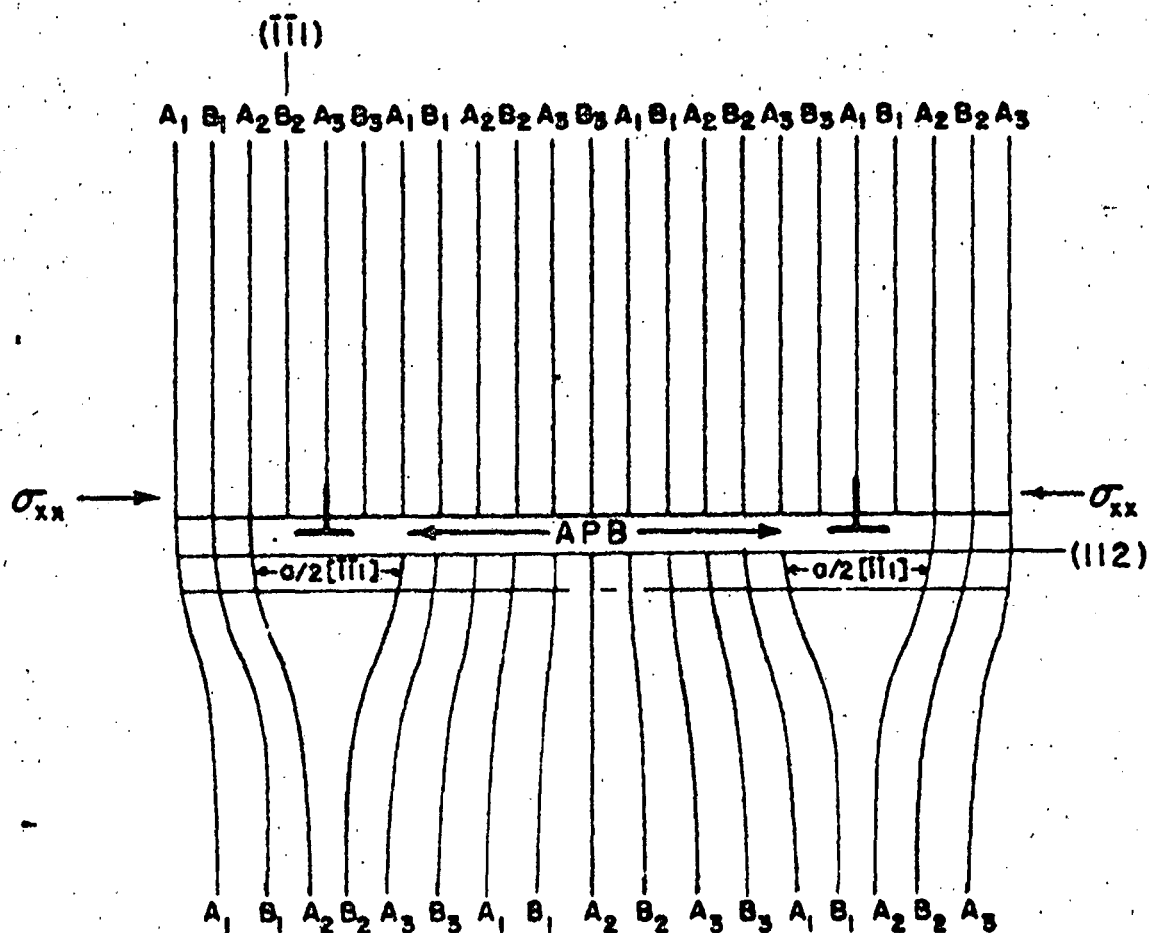


FIG. 8 DISSOCIATED SUPERDISLOCATION  
 $a [111] \rightarrow \frac{a}{2} [111] + \frac{a}{2} [111]$  ON THE  
 (112) PLANE. PLANE OF DRAWING IS (101).

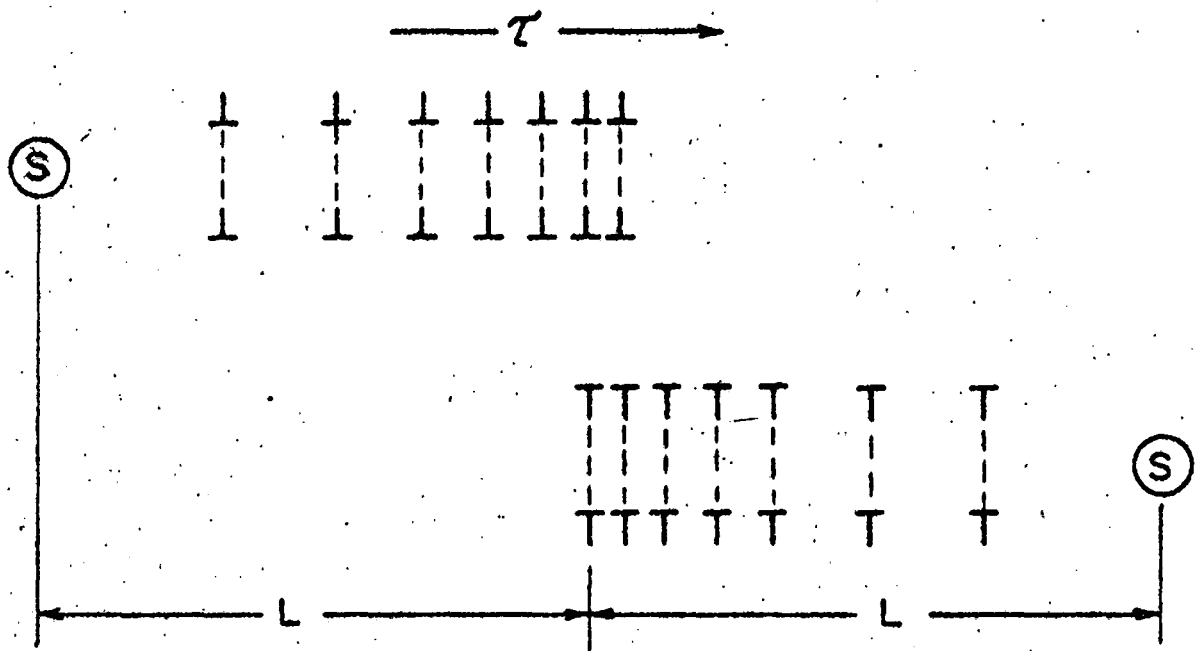


FIG. 9 EQUILIBRIUM CONFIGURATION OF  
SUPERDISLOCATION ARRAYS AT HIGH  
TEMPERATURES.

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