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# ON THE MECHANISM OF CREEP IN THE INTERMETALLIC COMPOUND AgMg AT ELEVATED TEMPERATURES

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June 24, 1964

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# ON THE MECHANISM OF CREEP IN THE INTERMETALLIC COMPOUND AgMg AT ELEVATED TEMPERATURES

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## ABSTRACT

The effects of temperature and stress on the creep rate were investigated from 733°K to 828°K in single crystals of the CsCl type of intermetallic compound AgMg oriented for slip in $\{\bar{3}21\}$ $<111>$ system. The activation energy for creep and the preexponential coefficient in the equation for creep rate agrees very well with a model where the rate controlling mechanism is one of viscous glide as modified by the presence of antiphase boundaries.

## I. INTRODUCTION

Recently Mukherjee and Dorn[1] demonstrated that $\{\bar{3}21\}$ <111> glide in single crystals of the long-range ordered intermetallic compound AgMg (CsCl type of lattice) obeyed the thermally activated Peierls mechanism over the range from 4° to 250°K (for the shear strain rate of $\dot{\gamma} = 4.16 \times 10^{-6}$) and obeyed an athermal mechanism from 250 up to somewhat above 350°K. The critical resolved shear stress for slip was observed to decrease precipitously with additional increases in temperature demonstrating that a second thermally activated mechanism becomes operative in this compound at yet higher temperatures. Because only limited information[2-4] is currently available on the high temperature mechanical behavior of long-range ordered alloys, it was deemed advisable to attempt an identification of the dislocation mechanisms responsible for creep in AgMg. In the context of this report we will demonstrate that creep in single crystals of AgMg by $\{\bar{3}21\}$ <111> glide from 733° to 828°K is most probably controlled by the rate of viscous creep as modified by the presence of antiphase boundaries.

## II. EXPERIMENTAL TECHNIQUE AND RESULTS

Oriented single crystal bars (3/8" in diam) of AgMg were grown under a He atmosphere by a modified Bridgman technique from chill cast master ingots made with Ag and Al each having purities better than 99.995 wt.%. The composition of the single crystals varied from 81.9 to 82.2 wt.% Ag and had therefore slightly higher Ag content than the stoichiometric amount of 81.6 wt.% for AgMg. Specimens were so oriented that the Schmid angles for $\{12\bar{3}\}$ <111> slip were $90-\chi_o \simeq \lambda_o \simeq 45 \pm 1°$.

The specimens were tested in a special creep unit so designed as to keep the resolved shear stress on the $\{12\bar{3}\}$ plane in the <111> direction constant to ±0.2%. Temperatures, as measured by thermocouples attached to the upper and lower grip sections of each specimen, were held constant to ±0.5°K of the reported values. To control oxidation of the specimens, creep was conducted under a positive atmosphere of He. Displacements of the strain gage, a linear differential transformer, were recorded autographically; the strains were measured to ±0.0001.

Over the range of temperatures and stress that were investigated, each specimen exhibited the usual creep curve, having a short transient stage, a long steady-state creep range and a tertiary stage leading to fracture; the latter was not studied in detail in the present investigation. A typical creep curve is shown in Fig. 1. Creep deformation in the specimens was observed to have taken place by $\{321\}$ <111> slip. After a steady-state creep rate was reached, further changes in the stress resulted in only very minor transients. Consequently, the data reported here were limited to the secondary stage of creep. For

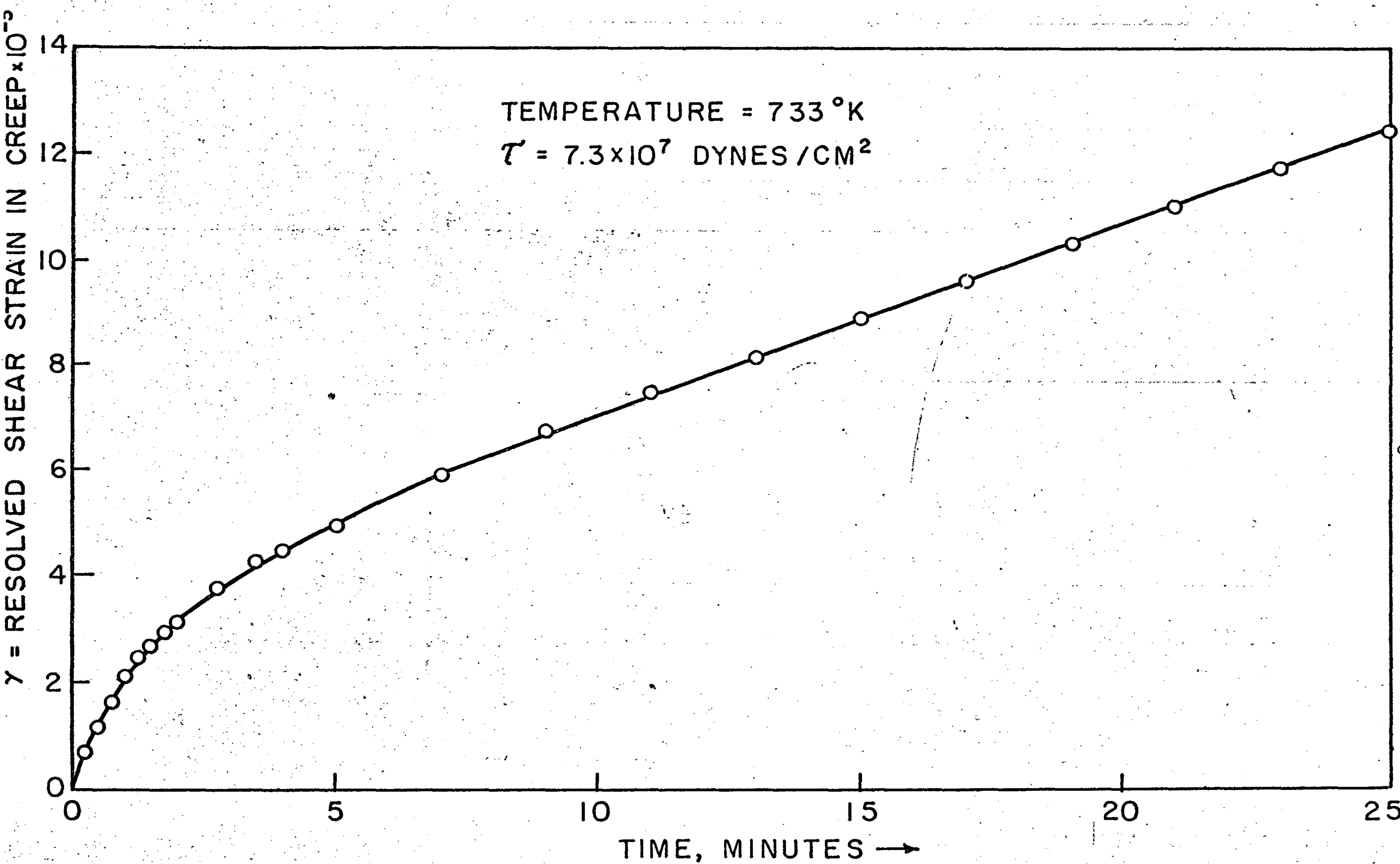


FIG. I. TYPICAL CREEP CURVE.

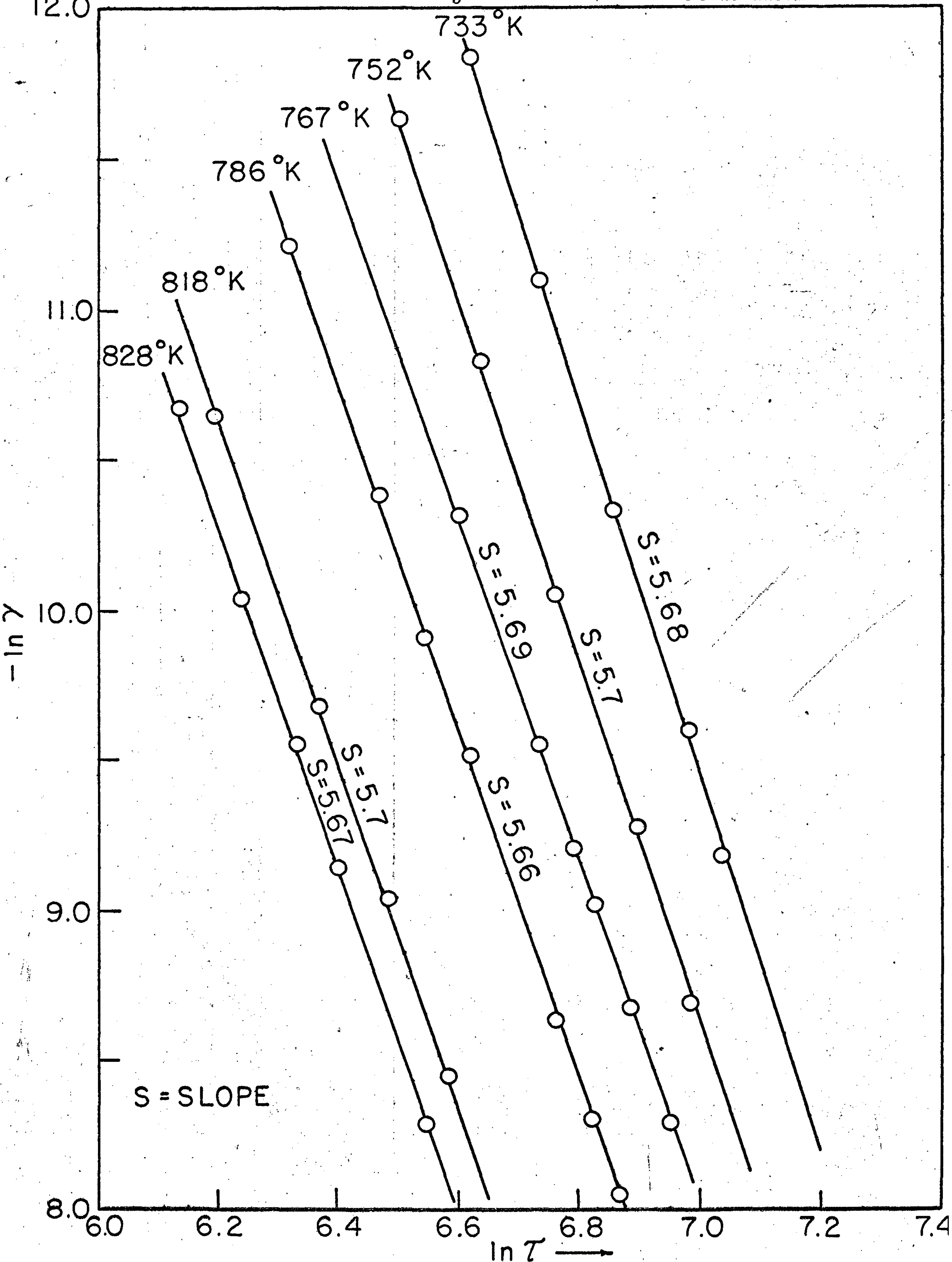
each test temperature, the stress was increased and decreased periodically and the resulting secondary creep rates were recorded as a function of the stress as shown in Fig. 2. These data were found to be independent of the order of increasing or decreasing the stress. Therefore, it was observed that the resolved shear strain rate, $\dot{\gamma}$, is given by

$$\dot{\gamma} \propto \tau^{n} \quad \text{where } n = 5.7 \pm 0.1 \tag{1}$$

where $\tau$ is the resolved shear stress on the slip plane and the exponent n is independent of the test temperature. Since this suggests that the apparent activation energy for creep is independent of the applied stress, the values of $\ell n(\dot{\gamma}/\tau^{n})$ were plotted in Fig. 3 as a function of 1/RT which reveals that

$$\dot{\gamma} = A'\tau^{n} e^{-\frac{Q'}{RT}} \tag{2}$$

where the apparent activation energy is given by $Q' = 47.5° \pm 2.0$kcal/mole, n = 5.7 and the coefficient $A' = 10^{-7.8}$cm$^{2}$/(dyne hour).


733°K
752°K
767°K
786°K
818°K
828°K
S=5.68
S=5.7
S=5.69
S=5.66
S=5.7
S=5.67
S = SLOPE
12.0
11.0
10.0
9.0
8.0
6.0
6.2
6.4
6.6
6.8
7.0
7.2
7.4
$-\ln\dot{\gamma}$
$\ln\tau$ ⟶

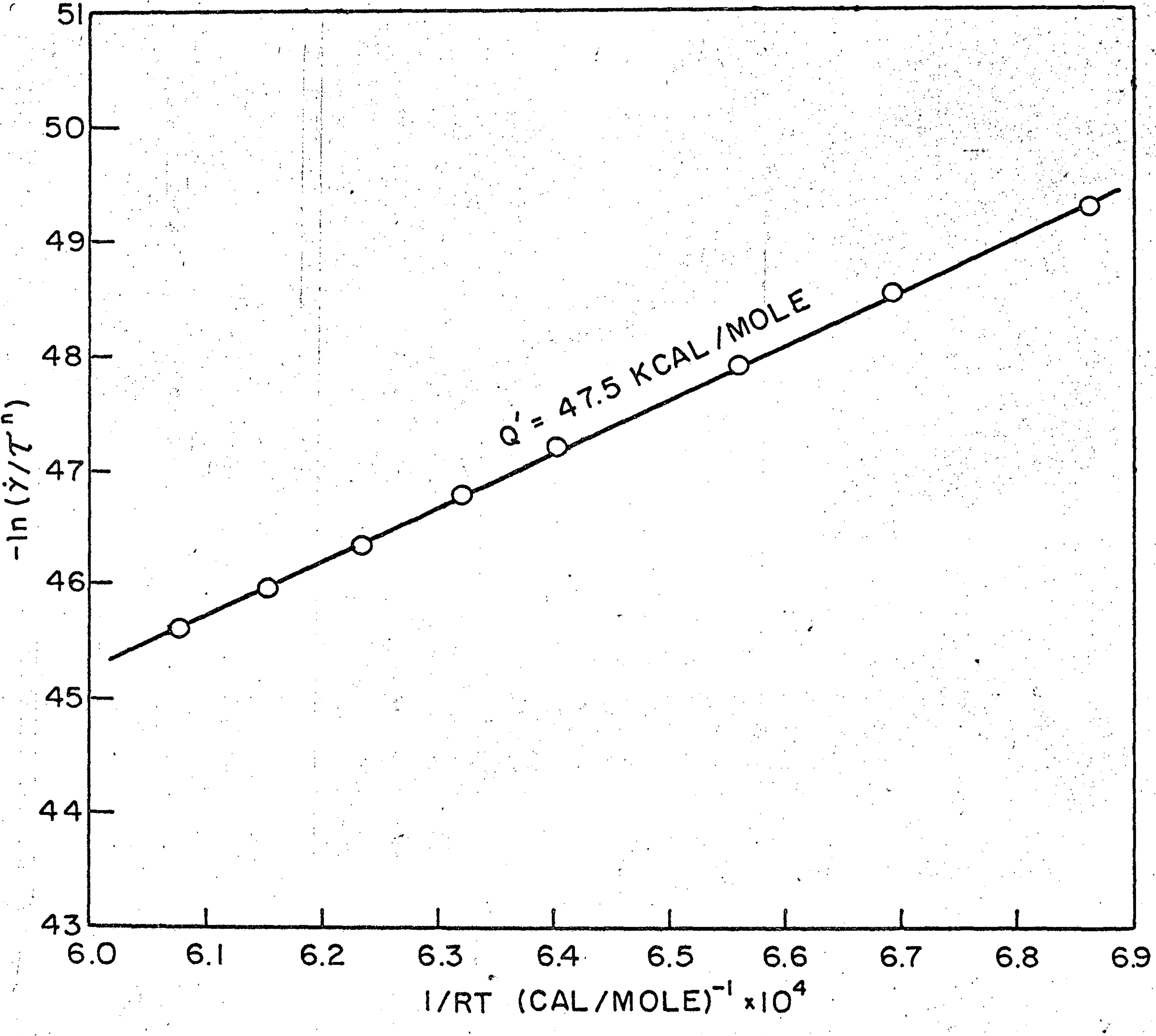

51
50
49
48
47
46
45
44
43
$-\ln(\dot{\gamma}/\tau^n)$
6.0
6.1
6.2
6.3
6.4
6.5
6.6
6.7
6.8
6.9
1/RT (CAL/MOLE)$^{-1}$ ×10$^4$
Q' = 47.5 KCAL/MOLE

FIG. 3.

## III. DISCUSSION

Because the apparent activation energy is insensitive to the applied shear stress on the slip plane, those dislocation mechanisms having activation energies that depend sensitively on the stress, such as the Peierls process,[5] the intersection mechanism,[6] motion of jogged screw dislocations,[7] dynamic recovery by cross slip,[8] etc. can be summarily disqualified as not being responsible for the creep of AgMg. The alternate mechanisms of either climb[9] of dislocations or viscous glide[10] of dislocation arrays, however, are compatable with the general form of the experimentally determined relationship given by Eq. 2. As will be elaborated later, however, both viscous glide and climb are influenced by the antiphase boundary energies in the long-range ordered AgMg alloy. To facilitate analyses of the creep data, we first estimate from thermodynamic data the antiphase boundary energy.

As shown by Flinn[2] the energy, σ per unit area of an antiphase boundary in the CsCl type lattice is given by

$$\sigma = \frac{-4hv}{a^2\sqrt{h^2 + k^2 + l^2}} \qquad h \geq k \geq l \tag{3}$$

where (h k l) is the plane of the antiphase boundary, a is the lattice constant and v is the enthalpy change for the bond equation:

$$-\frac{1}{2}(A-A) - \frac{1}{2}(B-B) + (A-B) = 0; \;\; v = v_{AB} - \frac{1}{2}(v_{AA} + v_{BB}) \tag{4}$$

where, v is negative. The value of v can be deduced from the enthalpy of formation of AgMg which has been shown by Hultgren et. al.[11] to be

$$-\mathrm{Ag}\left\{\begin{matrix}\mathrm{FCC}\\ \\ r_{Ag}=1.44\mathring{\mathrm{A}}\end{matrix}\right\}-\mathrm{Mg}\left\{\begin{matrix}\mathrm{CPH}\\ \\ r_{Mg}=1.60\mathring{\mathrm{A}}\end{matrix}\right\}+\mathrm{AgMg}\left\{\begin{matrix}\mathrm{CsCl}\\ \\ \frac{r_{Ag}+r_{Mg}}{2}=1.43\mathring{\mathrm{A}}\end{matrix}\right\}=0$$

$$\Delta H_1 = -10,400\mathrm{cal/mole} \qquad (5)$$

To obtain Eq. 4 we suggest that

$$-\mathrm{Ag}\left\{\begin{matrix}\mathrm{FCC}\\ \\ r_{Ag}=1.44\mathring{\mathrm{A}}\end{matrix}\right\}+\mathrm{Ag}\left\{\begin{matrix}\mathrm{BCC}\\ \\ r_{Ag}=1.38\mathring{\mathrm{A}}\end{matrix}\right\}=0 \quad \Delta H_2 = +289\mathrm{cal/mole} \qquad (6)$$

and

$$-\mathrm{Mg}\left\{\begin{matrix}\mathrm{HCP}\\ \\ r_{Mg}=1.60\mathring{\mathrm{A}}\end{matrix}\right\}+\mathrm{Mg}\left\{\begin{matrix}\mathrm{BCC}\\ \\ r_{Mg}=1.53\mathring{\mathrm{A}}\end{matrix}\right\}=0 \quad \Delta H_3 = 214\mathrm{cal/mole} \qquad (7)$$

where the enthalpy change from a close-packed to a BCC structure is taken to be about 0.1 of the latent hear of fusion as suggested by the data of Oelson et. al.;[12] the atomic radii, $r_{Ag}$ and $r_{Mg}$ are assumed to decrease 4% upon transformation in accord with Goldschmidt's rule. Subtracting Eqs. 6 and 7 from Eq. 5 gives

$$-\mathrm{Ag}\left\{\begin{matrix}\mathrm{BCC}\\ \\ r_{Ag}=1.38\mathring{\mathrm{A}}\end{matrix}\right\}-\mathrm{Mg}\left\{\begin{matrix}\mathrm{BCC}\\ \\ r_{Mg}=1.53\mathring{\mathrm{A}}\end{matrix}\right\}+\mathrm{AgMg}\left\{\begin{matrix}\mathrm{BCC}\\ \\ \frac{r_{Mg}+r_{Ag}}{2}=1.43\mathring{\mathrm{A}}\end{matrix}\right\}=0$$

$$\Delta H_4 = -10,900\ \mathrm{cal/mole} \qquad (8)$$

The strain energy associated with the change in $(r_{Ag} + r_{Mg})1/2$ for this reaction can be shown to be very small and is therefore neglected. On this basis

$$\frac{2NZv}{2} = \Delta H_4 \qquad (9)$$

where N is Avagadro's number and Z is the coordination number. Therefore Nv = -1362.5cal/mole of bond or $-0.9468 \times 10^{-13}$erg/mole of bond and, according to Eq. 3, $\sigma_{(321)} = 278\text{ergs/mole}^{-1}\text{cm}^{-2}$ for the antiphase boundary on the (321) plane and $\sigma_{(111)} = 200\text{ergs/mole}^{-1}\text{cm}^{-2}$ for the antiphase boundary on the (111) plane.

The separation of the dislocation pair comprising a superdislocation is given by

$$d = \frac{Gb^2}{2\pi\sigma}\left\{\cos^2\theta + \frac{\sin^2\theta}{1-\mu}\right\} \qquad (10)$$

where b is the Burgers vector of each dislocation of the dissociated superdislocation, namely b = a/2[111] where a is the lattice constant, G is the shear modulus of elasticity, $\mu$ is Poisson's ratio and $\theta$ is the angle between the Burgers vector and the dislocation line. Assuming that $G \simeq 18.0 \times 10^{10}\text{dynes/cm}^2$, $\mu \simeq 1/3$, we estimate that $5.89 \lesssim d/b \lesssim 8.84$ for the (321) plane, revealing that the dislocation pair of the dissociated superdislocation is appreciably separated. As shown by Cottrell[13] the athermal yield stress $\tau_B$ needed to produce antiphase boundaries as the dislocations traverse the antiphase domain boundaries is given by

$$\tau_B = \tau - \tau_G = \frac{\sigma}{L}\left(1 - \frac{\alpha a}{L}\right) \qquad (11)$$

where $\tau$ is the yield strength, $\tau_G$ is the athermal stress required to overcome the long-range interactions between dislocations, L is the mean diameter of an antiphase domain and $\alpha$ is a domain shape factor estimated to be about 6. Assuming that the previously determined athermal flow stress at intermediate temperatures arises almost exclusively from $\tau_B$ (i.e. $\tau_G$ is small) we obtain the reasonable lower bound value of L = 75Å cm for the mean antiphase domain diameter.

Flinn has shown that edge dislocations of a dissociated superdislocation will climb to different slip planes, at temperatures above about one-half of the melting temperature where diffusion is sufficiently rapid, as a result of the variation of the antiphase boundary energy with the plane of the boundary. Under these circumstances glide on a slip plane requires the formation of an antiphase boundary and the climb of a dissociated superdislocation also involves the formation of an antiphase boundary. When glide of a dislocation must produce an antiphase boundary without the aid of a thermal fluctuation, the athermal yield stress $\tau_A$ is given by

$$\tau_A b = (\tau_B + \tau_G) b = \sigma + \tau_G b \qquad (12)$$

If the applied stress is less than this value, glide can only take place when disordering occurs in advance of the dislocation as facilitated by diffusion. For the present we will consider Flinn's approach to this problem, applying small adjustments to his formulation of the theory.

Disordering in advance of the dislocation occurs as a result of diffusion and has about the frequency

$$\Gamma_+ = \nu Z e^{-\frac{(\Delta f + \mathcal{E}_v)}{kT}} e^{\frac{\tau b A_v}{kT}} \qquad (13)$$

where $\nu \simeq$ Debye frequency, $\Delta f$ is the free energy of activation for diffusion, $\mathcal{E}_v$ is the antiphase boundary energy for atomic area of the antiphase boundary on the slip plane, $\tau$ is the local stress in the direction of b, $A_v$ is the area per atom of the antiphase boundary, and kT has its usual meaning of the Boltzmann constant times the absolute temperature. The numerator of the exponent in the last term of Eq. 13 refers to the work done by the local stress, $\tau$, in assisting the formation of the antiphase boundary. Whereas Eq. 13 gives the frequency of the forward motion of the dislocation, disordering behind the dislocation causes a reversed motion with a frequency

$$\Gamma_- = \nu Z e^{\frac{-\Delta f + \mathcal{E}_v}{kT}} e^{-\frac{\tau A_v b}{kT}} \tag{14}$$

Consequently, the net velocity, v, of the forward motion of the dislocation is

$$v = b(\Gamma_+ - \Gamma_-) = 2\nu Z b e^{-\frac{(\Delta f + \mathcal{E}_v)}{kT}} \left\{\frac{\tau b A_v}{kT}\right\} \tag{15}$$

where the last term in the brackets refers to the degeneracy of the sinh $(\tau b A_v / kT)$ where $\tau b A_v < kT$. Since the diffusivity is given by

$$D = D_o e^{-\frac{\Delta h}{kT}} \simeq b^2 \nu e^{\frac{\Delta s}{k}} e^{-\frac{\Delta h}{kT}} \tag{16}$$

$$v = \frac{2Z\tau A_v D_o}{kT} e^{-\frac{(\Delta h + \mathcal{E}_v)}{kT}} \tag{17}$$

and the dislocation mobility B reduces to

$$B = v/\tau b = \frac{2ZA_v D_o}{bkT} e^{-\left(\frac{\Delta h + \mathcal{E}_v}{kT}\right)} \tag{18}$$

Weertman[10] has shown that at high temperatures under conditions where dislocation climb is more facile than glide, creep is controlled by viscous glide where

$$\dot{\gamma}_v \simeq \frac{2\pi(1-\mu)\,\tau^3 B}{G^2} \qquad (19)$$

Therefore when viscous creep controls the strain rate

$$\dot{\gamma}_v = K_v \, \frac{\tau^3}{G^2 T} \, e^{-\left(\frac{\Delta h + \mathcal{E}_v}{kT}\right)} \qquad (20)$$

where

$$K_v = \frac{4\pi Z A_o D_o (1-\mu)}{bk} \qquad (21)$$

If, however, viscous creep is more facile than creep by climb of dislocations, the creep rate will be controlled by climb. In this event, the creep rate can be approximated by Weertman's model, namely

$$\dot{\gamma}_c \simeq \frac{3\sqrt{6}\,\pi^2 \nu Z b^{1.5} (1-\mu)^{3.5} \tau^{4.5} e^{\frac{\Delta s}{k}} e^{-\frac{(\Delta h + \mathcal{E}_c)}{kT}}}{M^{1/2} G^{3.5} kT} \qquad (22)$$

where M is the number of dislocation sources per unit volume. Equation 22 differs from Weertman's[9] original equation by the term $\mathcal{E}_c$, the energy per atomic area of the antiphase boundary on the climb plane; this term accounts for the fact that as each dislocation of the dissociated dislocation climbs, it produces an antiphase boundary. When Eq. 16 is introduced, Eq. 22 reduces to

$$\dot{\gamma}_c = K_c \, \frac{\tau^{4.5}}{G^{3.5} T} \, e^{-\left(\frac{\Delta h + \mathcal{E}_c}{kT}\right)} \qquad (23)$$

where

$$K_c = \frac{3\sqrt{6}\,\pi^2 Z D_o (1-\mu)^{3.5}}{b^{1/2}\, k\, M^{1/2}} \qquad (24)$$

The three significant issues in arriving at a decision of the rate controlling mechanism for creep of AgMg concerns the activation energy, the stress law, and the coefficient K. Comparing Eqs. 23 and 20 with the experimentally determined creep rate reveals that the stress laws differ slightly. The theoretical expressions suggest therefore that a suitable activation energy can be estimated by suggesting the empirical relationship

$$\dot{\gamma} = K \frac{\tau^{5.7}}{G^{4.7}T} e^{-\frac{Q}{kT}}$$

where G is assumed to vary with temperature as previously announced.[1] The experimental data plotted in accord with this assumption, as shown in Fig. 4 reveal that $Q \simeq 50,800 \pm 800$cal/mole. To compare with theory we let $\Delta H = 40,600 \pm 900$cal/mole which is the activation energy for diffusion of $Ag^{110}$ in AgMg reported by Hagel and Westbrook.[14] Furthermore

$$\varepsilon_v = \sigma_{(123)}A_v \quad \text{and} \quad \varepsilon_c = \sigma_{(111)}A_c$$

where $A_c$ is the area per atom on the climb plane, which by simple calculation can be shown to be the (111) plane. Since $A = 1/2\, a^2\sqrt{h^2 + k^2 + l^2}$ for BCC crystals, $\varepsilon_v = -6v$ and $\varepsilon_c = -2v$. Using the previously determined value for v we find the theoretical activation energies to be

$$Q_v = 40,600 + 8,175 = 48,800 \text{ cal/mole}$$

$$Q_c = 40,600 + 2,725 = 43,300 \text{ cal/mole}$$

Since the calculated activation energy for viscous creep is in good agreement with the experimentally determined value, whereas that for climb is much too low, we conclude that the viscous creep mechanism controls the creep of AgMg. Finally, we calculate the theoretical

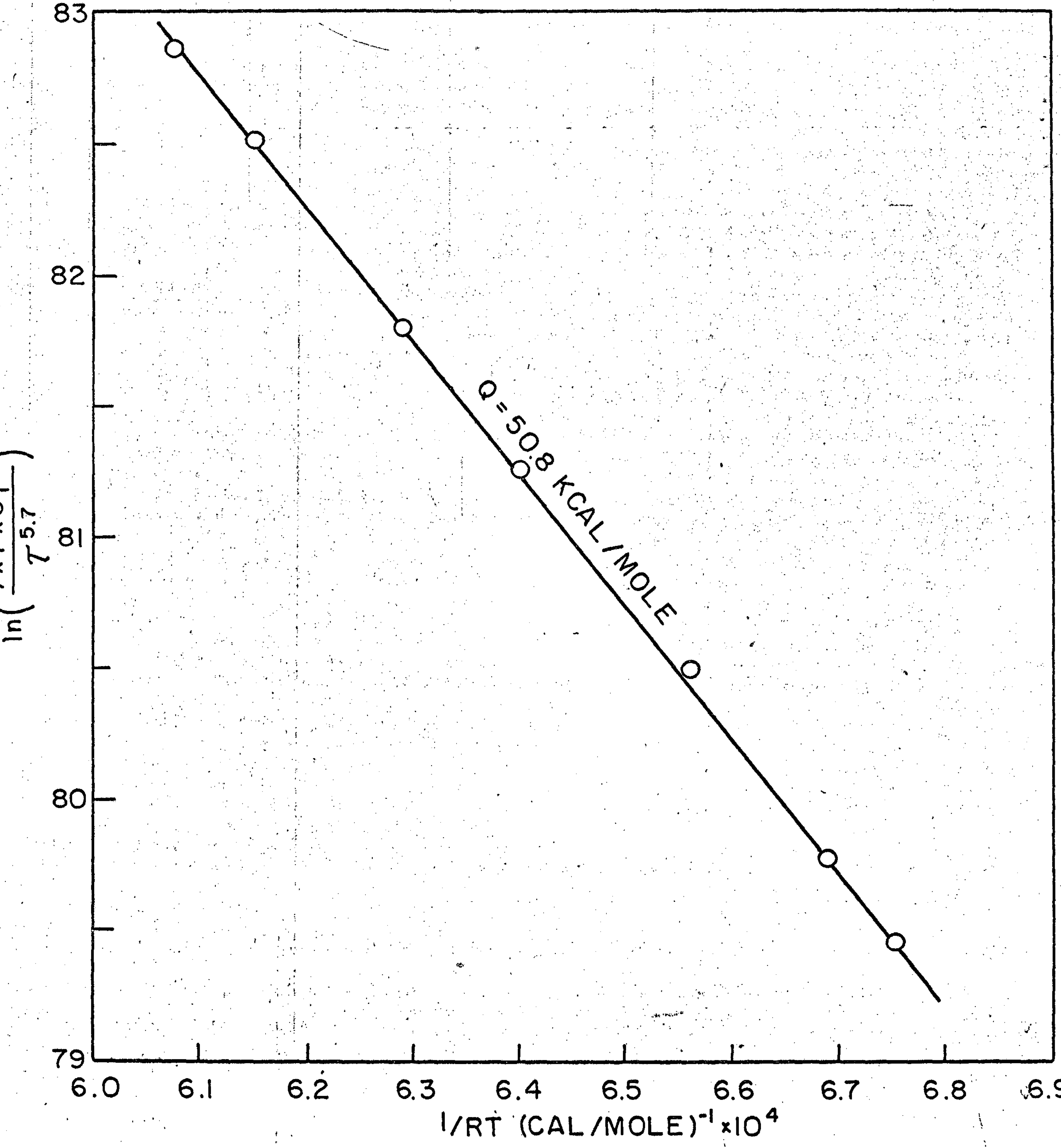

83
82
81
80
79
$\ln\left(\frac{\dot{\gamma} \times T \times G_T^{4.7}}{\tau^{5.7}}\right)$
Q = 50.8 KCAL/MOLE
6.0
6.1
6.2
6.3
6.4
6.5
6.6
6.7
6.8
6.9
$1/RT$ (CAL/MOLE)$^{-1}$ ×$10^4$

FIG. 4.

values of $\dot{\gamma}_v$ and $\dot{\gamma}_c$ for the mean values of $\tau = 5.0676 \times 10^7$ dynes/cm$^2$ and T = 818°K from which we obtained $\dot{\gamma}_v = 1 \times 10^{-4.7} < \dot{\gamma}_c = 1 \times 10^{8.6}$ which reveals that theoretically also viscous creep is the slower process and is therefore rate controlling.

Assuming appropriate values in the Eq. 21, the constant $K_v$ is found to be equal to $1 \times 10^{10.63} \frac{\mathrm{cm}^2}{\mathrm{dyne}}$ °K. A similar calculation for the value of the constant $K_c$ in Eq. 24 gives a value $K_c = 6.7 \times 10^{27} \frac{\mathrm{cm}^2}{\mathrm{dyne}}$ °K. The experimental values of creep rate $\dot{\gamma}$, creep stress $\tau$ and activation energy for creep, 50.8K cal/mole, when introduced in Eq. 20 gives a value of K(experimental) = $1 \times 10^{10.9} \frac{\mathrm{cm}^2}{\mathrm{dyne}}$ °K, which again suggests that the viscous creep process is rate controlling.

The observed value of the stress dependence of creep rate $n \simeq 5.7$ is higher than that predicted from the theory of viscous creep. This tendency of higher values of n is also noticed in some other instances of viscous creep and the reason for this may lie in the details of the model used in the theory.

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