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Authors
Murty, K. Linga
Mohamed, F.A.
Dorn, J.E. (LBL).

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EFFECT OF VACANCY SINKS AND SOURCES ON SERRATED YIELDING DUE TO SOLUTE LOCKING

K. Linga Murty*, F. A. Mohamed** and J. E. Dorn†
Inorganic Materials Research Division
Lawrence Berkeley Laboratory
University of California, Berkeley, California

It is generally accepted that a realistic model for the Portevin-Le Chatelier effect due to excess vacancy stimulated solute diffusion and dislocation locking must consider vacancy annihilation as well as vacancy creation during straining (1). Modified versions (2,3) of Cottrell model (4) of repeated yielding take into account the thermal vacancy contribution as well as the strain-produced vacancies but fail to consider annealing of vacancies to appropriate sinks. Recent experimental data on Al-Mg alloy by MacEwen and Ramaswami (3) reveal important deviations from Cottrell model which they qualitatively accounted for by vacancy annihilation. The present note is an attempt to quantitatively explain these deviations by taking into account the disappearance of vacancies at sinks.

The Cottrell model has been most used and abused in attempts to explain the experimental findings of the delay of plastic strain ($\varepsilon_0$) and critical strain-rate ($\dot{\varepsilon}_c$) for the appearance of serrations and their temperature dependencies. According to Cottrell (4) the critical strain-

*Research Associate
**Research Assistant
†Senior Scientist and Professor of Materials Science respectively of the College of Engineering, University of California, Berkeley, California; deceased September, 1971
rate at which serrations start appearing is given to be

\[ \dot{\epsilon}_c = K \rho_m C_v e^{-E_m/RT}, \]  

where \( K \) is a constant, \( \rho_m \) the mobile dislocation density, \( C_v \) the vacancy concentration and \( E_m \) the activation energy for vacancy migration. Except for two cases (2,3) \( C_v \) in the above equation was taken to be the total concentration of vacancies produced during straining whereas in strict sense \( C_v \) is the net total vacancy concentration and

\[ C_v = C_v^T + C_v^e - C_v^- , \]

Where \( C_v^T \) is the thermal equilibrium vacancy concentration at the test temperature \( T \), \( C_v^e \) the total excess vacancy concentration produced by straining at strain \( \epsilon \) and \( C_v^- \) the concentration of annihilated vacancies at temperature \( T \), and \( C_v^- \leq C_v^e \).

Point defects are generated by the non-conservative motion of jogs on screw dislocations (5) and the defects thus produced are vacancies or interstitials depending upon the nature of the jogs. Assuming that vacancies are created at vacancy producing jogs on gliding screws then the rate of production of these excess vacancies is given by

\[ \dot{C}_v^e = \frac{\rho_s v_s p}{N l_j b}, \]

where \( \rho_s \) is the density of mobile screw dislocations, \( v_s \) the velocity of these screw dislocations, \( p \) the height of jogs in Burgers vectors, \( l_j \) the jog separation, \( b \) the Burgers vector and \( N \) the number of lattice sites per unit volume. Using \( \dot{\gamma} = \frac{3}{4} \dot{\epsilon} = \rho_s b v_s \) we find (6)

\[ \dot{C}_v^e = \frac{3}{4N} \left( \frac{p}{l_j b^2} \right) \dot{\epsilon}, \]
where \( \dot{\varepsilon} \) is the imposed tensile strain-rate. Thus if there is no recovery, i.e. once formed a vacancy does not anneal out to a sink, we find from Eq. 3 that

\[
C_v^\varepsilon = \frac{3}{4N} \left( \frac{p}{1_j b^2} \right) \varepsilon ,
\]

[4]

In reality however such excess vacancies disappear at sinks at a rate proportional to the number produced so that

\[
\dot{C}_v = \left( \frac{4D_v}{L_s^2} \right) C_v^\varepsilon ,
\]

[5]

where \( D_v \) is the vacancy diffusivity and \( L_s \) the vacancy source-sink distance. In a very fine grained material this \( L_s \) may be taken as the grain size, but in general dislocations are the primary vacancy sinks in which case

\[
L_s^2 = \frac{1}{\rho} ,
\]

[6]

so that

\[
\dot{C}_v = 4D_v \rho K_1 \varepsilon ,
\]

[7]

with \( K_1 = \frac{3}{4N} \frac{p}{1_j b^2} \). It has been well established (7) that the density of dislocations, \( \rho \), varies with the square of the flow stress and \( \sigma = \alpha Gb\sqrt{\rho} \) with \( \alpha \approx 1 \). On the other hand \( \sigma \approx \sigma (\varepsilon) \) dependent on conditions and in general \( \sigma = A \sqrt{\varepsilon} \) so that \( \rho = \rho_0 \varepsilon \) while in single crystals for Stage II deformation \( \sigma = B (\varepsilon - \varepsilon_0) \) so that \( \rho \sim \varepsilon^2 \). Substituting the explicit strain-dependence of the dislocation density and integrating we obtain

\[
\dot{C}_v = \frac{4 D_v K_1}{3 \dot{\varepsilon}} \rho \varepsilon^2
\]

[8]
A similar but more complicated expression may be obtained for single crystals. Thus the net total vacancy concentration present at strain $\varepsilon$ and temperature $T$ is

$$C_v = C_v^T + C_v^e - C_v^{-} = e^{-\left(E_f + TS_f\right)/RT} + \frac{3p}{4N_1b^2} \left( \varepsilon - \frac{4D_v\rho}{3\bar{\varepsilon}} \right)^2$$

where $E_f$ and $S_f$ are the activation energy and entropy respectively of vacancy formation.

As an illustration of the applicability of the above analysis recent data of MacEwen and Ramaswami (3) on serrated yielding in Al-Mg alloy will be considered. Fig. 1 is a reproduction of their data on the temperature dependence of $\varepsilon_o$. The datum point in this figure at $\varepsilon_o = 1.23 \times 10^{-3}$ falls above the straight line obtained from other points based on Cottrell equation. If the vacancy annihilation at the corresponding temperature ($370^\circ K$) were negligible the datum point would have been at $\varepsilon_{\text{new}} = 5.0 \times 10^{-4}$. Because some vacancies annealed out it needed more than this $\varepsilon_{\text{new}}$. Hence

$$C_v^e \{\varepsilon_o\} - C_v^{-} \{\varepsilon_o\} = C_v^e \{\varepsilon_{\text{new}}\}$$.  

Or, from Eqs. [4] and [9]

$$C_v^e \{\varepsilon_{\text{new}}\} = K_1 \varepsilon_{\text{new}} = K_1 \left[ \varepsilon_o - \frac{4D_v\rho}{3\bar{\varepsilon}} \right]^2$$

Inserting the experimental values of $\bar{\varepsilon}$, $\varepsilon_o$, $D_v$ at $370^\circ K$ and the extrapolated value of $\varepsilon_{\text{new}}$, a value of $2 \times 10^9 \text{ cm}^{-2}$ is obtained for the dislocation density. The value for $\rho$ is of the same order (indeed 3 times) as expected from Ham-Jaffrey (8) equation,

$$\rho = 4.73 \times 10^{11} \varepsilon = 6 \times 10^8$$

Thus the present analysis yields excellent agreement with the experimental findgs. With the above modifications then the Cottrell equation for the critical (imposed) strain-rate may be rewritten to be

$$\dot{\varepsilon}_c = K' \varepsilon_o e^{-\frac{E_m}{RT}} \left\{ e^{-\frac{E_f}{RT}} + \frac{3p_e^{-}}{4N_1b^2} \left( 1 - \frac{4D_v\rho}{3\bar{\varepsilon}_c} \right) \right\}$$
in lieu of the simplified equation used thus far.

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References

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\[ \ln \varepsilon = 5.6 \times 10^{-5} \text{SEC}^{-1} \]

\[ \varepsilon_0 = 1.23 \times 10^{-5} \]

\[ \varepsilon_0 = 5 \times 10^{-4} \]

\[ (370^\circ \text{K}) \]

\[ 1000/ T, \, ^\circ \text{K}^{-1} \]

FIG. 1

Plot of \( \ln \varepsilon_0 \) vs \( 1/T \). Data of MacEwen and Ramaswami. (3)
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