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Void Lattice Formation
— Spinodal Decomposition of Vacancies

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Abstract

We consider the mechanism of the void lattice formation as a spinodal decomposition of vacancies. In the linear region of the spinodal decomposition, the void lattice constant is determined by estimating the periodicity growing in the spinodal decomposition. We assume that the void radius in the void lattice should be determined from the static stability according to Bullough and Stoneham et al..

§1. Introduction

In 1970 Evans¹⁾ discovered that superlattices of voids are formed under a certain condition in a metal irradiated by neutrons or heavy ions. Such a superlattice has been observed in various sorts of metallic crystals, including Mo^{2),3)}, Ta²⁾, Nb^{2),4)}, W⁵⁾, Ni⁶⁾, and Al⁷⁾, and also in certain alloys (eg. TZM⁸⁾, and Mo_{0.95}Ti_{0.05}²⁾).

A superlattice of voids has been found to be formed only under conditions of heavy radiation damage (> 1 dpa.) near the threshold temperature of void formation $\sim 0.3T_m$, where T_m is the melting temperature of the metal. The void superlattice has the same lattice structure as the host lattice, and the void lattice axes are also parallel to those of the host. For example in Mo, Ta, Nb, and W the void lattice forms a bcc array and in Ni and Al it forms a fcc array.

The distance between neighbouring voids in a lattice (i.e. void lattice constant) ranges from 200\AA to 800\AA , depending on the radiation conditions. Void sizes also have a wide range between 40 and 250\AA . The ratio of the void lattice constant to the void size falls in the range between 5 and 15.

A void superlattice usually has some defects as its host crystal does. For example, vacant sites, dislocations, variations of void size, void clusters consisting of several smaller voids on one site, and so on have been found.

Stability of the void lattice is also a characteristic feature. Once a void lattice is formed, coarsening of the void does not take place even in the annealing temperature ($\sim 1100^\circ\text{C}$ in case of Mo), and no growth of the void size occurs with further radiation. It is very different from the case of isolated voids, in which both the coarsening and the growth easily occur at the annealing temperature.

There is no evidence of voids moving in the process of the void lattice formation. Aside from the void lattice the ordering of voids into rows is observed in some cases.⁹⁾ It seems to be a "germ" of a void lattice.

So far there have been some theoretical treatments about the void lattice. Especially about the static stability of a void lattice, the elastic energy was calculated in various ways. These are reviewed briefly in the following.

The first calculation is due to Malen and Bullough¹⁰⁾. Assuming that voids already formed a lattice, they calculated the elastic energy as a function of a parameter L/R (R : the void size, L : the void lattice constant). They assumed an appropriate form for the force field around a void. Thus, using the anisotropic elastic Green's function they summed up the elastic fields originated from other voids. In this calculation they showed that the elastic energy takes a minimum value at a certain value of L/R . Stoneham¹¹⁾ has performed similar calculations in a Fourier transformed form. Let $F_\alpha^*(\mathbf{r})$ be the α -component of the force field ($\alpha = x, y, z$) arising from a void at the origin. He assumed that F_α^* is given by

$$F_{\alpha}^{*}(\mathbf{r}) = \sum_i F^{*} \delta(\mathbf{r} - \delta_i) \frac{\delta_{i\alpha}}{|\delta_i|} , \quad (1.1)$$

where $\{\delta_i\}$ are vectors which are suitably chosen depending on the shape of the void (provided that $|\delta_i|$ is of the same order as the void size.) In eq. (1.1) F^{*} is a constant. If $\{\mathbb{L}_j\}$ is the set of the co-ordinates of voids including the origin, a Fourier transformed form of the net force field at the origin is expressed as follows

$$F_{\alpha}(\mathcal{Q}) = \sum_j' e^{i\mathcal{Q} \cdot \mathbb{L}_j} \sum_{\mathbf{r}} e^{-i\mathcal{Q} \cdot \mathbf{r}} F_{\alpha}^{*}(\mathbf{r}) , \quad (1.2)$$

where the summation \sum_j' is taken over all voids except the one at the origin. Using the anisotropic Green's function, the elastic energy is given by

$$E = -\frac{1}{2N} \sum_{\mathcal{Q}} F_{\alpha}(\mathcal{Q}) G_{\alpha\beta}(\mathcal{Q}) F_{\beta}(-\mathcal{Q}) , \quad (1.3)$$

where N is the number of lattice sites per volume, and the summations of \mathcal{q} is taken within the first Brillouin zone of the host lattice, and only over the reciprocal lattice points of the void lattice because of the periodicity of the voids. By summing up these numerically Stoneham has shown that E takes a minimum value when $L/R \sim 3$.

Tewary and Bullough¹²⁾ considered the effects of voids on the elastic properties of the host lattice. According to them,

E has a minimum value at $L/R \sim 10$ which agrees with the experiments.

Since these calculations are based on the Green's function derived from the continuum elastic theory, they can only calculate the energy as a function of L/R . Thus the absolute value of L is not fixed by these static stability analysis. Therefore it is possible that some other dynamical mechanisms may determine L . In this paper we show that the spinodal decomposition of vacancies determines the value of L , and with the value of L thus obtained, R is determined so as to satisfy the static stability condition.

The free energy of the system including vacancies is presented in §2. In §3 the spinodal decomposition of vacancies is discussed in the linear region, and the growth of periodicity is discussed. Comparison with experimental conditions and conclusion are given in §4. A proposal for a new experiment is also discussed in §4.

§2. Free Energy

A crystal lattice under heavy radiation damage contains a great number of imperfections. Though there are various sorts of reactions and interactions between various kinds of imperfections, we consider here only vacancies. In order to describe an irradiated solid crystal which includes vacancies, we start from the following free energy per atom.

$$\begin{aligned}
 F = \frac{N}{N_0} & \left[\epsilon c(r) + \frac{1}{2} c(r) \sum_i c(r+R_i) U(R_i) \right. \\
 & + k_B T \{ c(r) \log c(r) + (1-c(r)) \log (1-c(r)) \\
 & \left. + \frac{N_0}{N} f_{\text{vib}}(c(r)) \} \right], \tag{2.1}
 \end{aligned}$$

where N and N_0 are the numbers of the lattice sites and the atoms per unit volume respectively, and $c(r)$ is the vacancy concentration. Note that N_0 is equal to $(1-c(r))N$. The formation energy of a vacancy is ϵ , and $U(r)$ is the interaction energy for a vacancy pair separated at the distance r . Lattice site vectors is represented by the set of $\{R_i\}$. The vibrational free energy f_{vib} will be discussed later. Formally we take a concentration expansion of f_{vib} as

$$f_{\text{vib}}(c(r)) = f_{\text{vib}}^{(0)} + f_{\text{vib}}^{(1)} c(r) + f_{\text{vib}}^{(2)} c(r)^2 + \dots \tag{2.2}$$

The vacancy-vacancy interaction in metal consists of the electronic part and the elastic part. If two vacancies are

at a nearest neighbour lattice sites, the electronic part dominates over the elastic part¹³⁾. However the electronic part decreases exponentially with increasing distance of separation. To estimate the elastic part^{14),15)} we must solve the problem consistently with the electronic states. Even if the effect of the conduction electrons is taken into account in a simplified form of the interatomic potential (for example, the Morse potential), we must evaluate relaxation of many atoms around a vacancy in the case of the large separation of two vacancies. But in general the elastic part decreases more quickly than r^{-2} .^{14),15)} Therefore we may take into account only the nearest neighbour site interaction. Because two vacancies placed at the nearest neighbour lattice sites are nothing but a divacancy, we have only to know the binding energy of a divacancy. There are some attempts to calculate¹³⁾ the binding energy of a divacancy for metals and some experiments¹⁶⁾ have also been done. They are generally about 0.1 ~ 0.3 eV for many sorts of metals (see Table I). For Mo or Nb etc. it might be a little more than 0.3 eV because of the large cohesive energy. The vacancy formation energy ϵ is about 0.5 ~ 1.0 eV for noble metals and Al, and is about 2 ~ 3 eV for bcc transition metals (Nb, Mo, and W) (see Table II).

When the concentration of vacancies is weakly nonuniform, we take the gradient expansion of the concentration as

$$C(\mathbf{r} + \mathbf{R}_i) = C(\mathbf{r}) + \mathbf{R}_i \cdot \nabla C(\mathbf{r}) + \frac{1}{2} (\mathbf{R}_i \cdot \nabla)^2 C(\mathbf{r}) + \dots, \quad (2.3)$$

and write eq. (2.1) as

$$\begin{aligned}
 F[C(\mathbf{r})] = & \varepsilon C(\mathbf{r}) + \alpha C(\mathbf{r})^2 + \frac{1}{6}(\alpha - \varepsilon)a^2 C(\mathbf{r}) \nabla^2 C(\mathbf{r}) \\
 & + k_B T (f_{\text{vib}}^{(1)} C(\mathbf{r}) + f_{\text{vib}}^{(2)} C(\mathbf{r})^2) \\
 & + k_B T \left\{ \frac{C(\mathbf{r}) \log C(\mathbf{r})}{1 - C(\mathbf{r})} + \log(1 - C(\mathbf{r})) \right\},
 \end{aligned}$$

$$\alpha = \varepsilon + \frac{1}{2} \sum_i U(\mathbf{R}_i), \quad (2.4)$$

where we set $a \equiv |\mathbf{R}_i|$ (only the nearest neighbour lattice site interaction is taken into account).

§3. Void Lattice Formation Process

The evolution of a system including vacancies is described by

$$\frac{\partial c}{\partial t} = D \nabla \cdot c(r) \{1 - c(r)\} \nabla \frac{\partial F}{\partial c(r)} + q, \quad (3.1)$$

where D is the diffusion constant of a vacancy, and q is the influx of vacancies from outside the system. A derivation of eq. (3.1) is given in Appendix. The directions of the growth of the inhomogeneity is parallel to the axes of the host crystal because of the existence of the anisotropic elastic energy and the growth in each direction of the host lattice axes is independent of the ones in other directions. From now on we will analyze the growth of $c(r)$ in a particular direction thus reducing the problem to a one-dimensional one.

We expand $c(x)$ around the mean concentration c_0 (e.g. $c(x) = c_0 + \gamma(x)$). If we assume that $1 \gg c_0 \gg \gamma$, eq. (3.1) is transformed into

$$\frac{\partial c}{\partial t} = - \frac{DA}{k_B T} \left[(C_0(t) - C_s) \frac{\partial^2 \gamma(x,t)}{\partial x^2} + C_0(t) \mu \frac{\partial^4 \gamma(x,t)}{\partial x^4} \right] + q, \quad (3.2)$$

$$C_s \equiv \frac{-k_B T}{2(\alpha + f_{vib}^{(2)} k_B T)},$$

$$A \equiv -2(\alpha + f_{vib}^{(2)} k_B T),$$

$$\mu \equiv \frac{1}{6} \frac{(\alpha - \epsilon) a^2}{\alpha + f_{vib}^{(2)} k_B T}.$$

In eq. (3.2) only the lowest terms in c_0 and γ are retained, and the gradient expansion is retained up to the fourth order. Taking the Fourier transformation of eq. (3.2), we obtain

$$\frac{\partial \gamma(k, t)}{\partial t} = \frac{DA}{k_B T} [(\varrho t - C_s) k^2 \gamma(k, t) - \varrho t \mu k^4 \gamma(k, t)]. \quad (3.3)$$

If we define t_0 as

$$t_0 = \frac{C_s}{\varrho(1 - \mu k^2)}, \quad \mu k^2 < 1, \quad (3.4)$$

the growth of $\gamma(k)$ occurs after the time t_0 . For $k^2 > \mu^{-1}$ the growth of γ never occurs.

Let us denote by γ_0 the initial deviation of the concentration from the uniform one, which is caused by the thermal fluctuation in the system. Then we obtain

$$\gamma(k, t) = \gamma_0 \exp \left[\int_{t_0}^t dt \frac{DA}{k_B T} \{ (\varrho t - C_s) k^2 - \mu \varrho t k^4 \} \right], \quad t > t_0. \quad (3.5)$$

The mode which gives the largest value of γ is given by

$$k_{\max}^2 = \frac{1}{2\mu} \left(1 - \frac{2C_s}{\varrho(t+t_0)} \right) \quad (3.6)$$

Thus the maximum value of $\gamma(k)$ among all modes is given by

$$\gamma(k_{\max}, t) = \gamma_0 \exp \left[\frac{2DA C_s^2 \mu^2}{k_B T \varrho} k_{\max}^6 \frac{1}{(1 - 2\mu k_{\max}^2)^2 (1 - \mu k_{\max}^2)} \right], \quad (3.7)$$

We show k_{\max} dependence of $\gamma(k_{\max})$ in Fig.1. In drawing this figure, it is necessary to fix the value of $f_{\text{vib}}^{(2)}$. Although there have been some attempts¹⁷⁾ to determine the value of $f_{\text{vib}}^{(2)}$, it is very difficult to do it precisely in a reliable way for any metal. But usually the value of $f_{\text{vib}}^{(2)} k_B T$ is considered to be of the same order of ϵ or $\sum_i U(R_i)$ at the temperature under consideration. Therefore within an order estimation we can set $\alpha + f_{\text{vib}}^{(2)} k_B T \sim -1\text{eV}$. In Fig.1, some reasonable values for the experimental conditions on the transition metal (Mo, Nb, and W) are assumed; namely $D=10^3 \text{a}^2 \text{sec}^{-1}$, $q=10^{-7} \text{sec}^{-1}$, $T=10^3 \text{K}$, $\epsilon=2\text{eV}$, $U=-0.4\text{eV}$ and the nearest neighbour lattice site number $z=8$ (for bcc lattice). We take values of ϵ , and U from existing literatures^{13), 16)}, but we choose D and q rather arbitrarily, because these values are not well known. Nevertheless, the result turns out to be quite insensitive to this choice; eq. (3.7) and Fig.1 show that a very sharp growth of γ occurs at about $ka \approx 10^{-1}$, and we find that the variations of q , D , and $f_{\text{vib}}^{(2)}$ does not so much affect the value of k_{\max} at the sharp rise of the curve in Fig.1. This is because of k_{\max}^6 dependence in the exponential part of eq. (3.7).

§4. Discussion

So far we found that the mechanism of the spinodal decomposition of vacancies gives the periodicity $\lambda=2\pi/k$ which agrees with the void lattice constant. Of course the above discussion is limited to the linear region (initial region) of the spinodal decomposition. It is, however, reasonable that the mode which is determined in the linear region grows even in the later stages, since the increase of γ is very sharp.

Since we don't know the precise data of q and D and so on in the experimental conditions, we cannot estimate the explicit value of the void lattice constant for a particular experiment. We can only say that the void lattice constant should have a tendency to become large for a larger diffusion constant. This assertion is confirmed by the fact that the void lattice constant has a larger value for Ni or Al than for Mo, Nb, or W in experiments; in fact transition metals have smaller diffusion constants than other metals.

It is necessary to measure the value of q and D in an experiment simultaneously as well as other properties of the void lattice. Then we shall be able to discuss the details and the quantitative propriety of eq. (3.7).

Appendix

A derivation of eq. (3.1) is given in this Appendix. The starting point is the kinetic Ising model, which is used by Saito¹⁸⁾.

We consider a binary alloy consisting of two components A and B. Of course in the case of the system including vacancies, we may regard that the component A is the vacancy and the component B is the atom. We divide the system into small cubic cells of a size ℓ in the atomic distance unit. Note that one cell contains ℓ^3 atoms. We assume that the size of a cell is large enough to define the concentration of a cell, and is small enough to be in a uniform density. If we define the concentration of A in the μ -th cell as n_μ ($0 \leq n_\mu \leq 1$), we can give the free energy as a function of the local concentration n_μ . This means

$$F = F(n_1, n_2, \dots, n_\mu, \dots, n_N)$$

If the concentrations of cells are given by $n_1, n_2, \dots, n_\nu, \dots, n_\kappa, \dots$ at first, the transition probability of an exchange between a B atom in the ν -th cell and an A atom in the κ -th cell is assumed to be

$$W = n_\kappa(1-n_\nu) \exp \left[\left\{ E(\dots, n_\nu, \dots, n_\kappa, \dots) - E(\dots, n_\nu + \frac{1}{a^3}, \dots, n_\kappa - \frac{1}{a^3}) \right\} / 2k_B T \right] \quad (\text{A.1})$$

Then the probability distribution $P(\{n_\mu\})$ for the configuration of the set of $\{n_\mu\}$ obeys

$$\begin{aligned} \frac{\partial P(\{n_{\mu}\})}{\partial t} = & -\frac{a^4}{2\tau} \sum_{\langle \nu\kappa \rangle} \left[n_{\kappa}(1-n_{\nu}) \exp\left[\frac{1}{2k_B T} \left\{ E(n_{\mu}, n_{\kappa}) \right. \right. \right. \\ & \left. \left. \left. - E\left(n_{\nu} + \frac{1}{a^3}, n_{\kappa} - \frac{1}{a^3}\right) \right\} \right] P(n_{\nu}, n_{\kappa}) \right. \\ & \left. - \left(n_{\nu} + \frac{1}{a^3}\right)(1-n_{\kappa} + \frac{1}{a^3}) \exp\left[\frac{-1}{2k_B T} \left\{ E(n_{\nu}, n_{\kappa}) \right. \right. \right. \\ & \left. \left. \left. - E\left(n_{\nu} + \frac{1}{a^3}, n_{\kappa} - \frac{1}{a^3}\right) \right\} \right] P\left(n_{\nu} + \frac{1}{a^3}, n_{\kappa} - \frac{1}{a^3}\right) \right] \quad (\text{A.2}) \end{aligned}$$

where τ is a rate constant of an exchange, and $\langle \nu\kappa \rangle$ denotes all the combinations of neighbouring cells. We assume that an exchange takes place only at the boundary surface of two neighbouring cells (this surface area is l^2). We determine the free energy so as to satisfy the detailed balance. This is expressed as

$$F(\{n_{\mu}\}) \equiv E(\{n_{\mu}\}) + k_B T \sum_{\{\mu\}} \left\{ n_{\mu} \log n_{\mu} - n_{\mu} + (1-n_{\mu}) \log(1-n_{\mu}) - (1-n_{\mu}) \right\} \quad (\text{A.3})$$

Using eq. (2.A.3) we can transform eq. (2.A.2) into

$$\begin{aligned} \frac{\partial P(\{n_{\mu}\})}{\partial t} = & -\frac{a^4}{2\tau} \sum_{\langle \nu\kappa \rangle} \left[\sqrt{n_{\kappa}(1-n_{\kappa} + \frac{1}{a^3})(n_{\nu} + \frac{1}{a^3})(1-n_{\nu})} \times \right. \\ & \left. \left\{ \exp\left[\left\{ F(n_{\nu}, n_{\kappa}) - F\left(n_{\nu} + \frac{1}{a^3}, n_{\kappa} - \frac{1}{a^3}\right) \right\} / 2k_B T \right] P(n_{\nu}, n_{\kappa}) \right. \right. \\ & \left. \left. - \exp\left[-\left\{ F(n_{\nu}, n_{\kappa}) - F\left(n_{\nu} + \frac{1}{a^3}, n_{\kappa} - \frac{1}{a^3}\right) \right\} / 2k_B T \right] P\left(n_{\nu} + \frac{1}{a^3}, n_{\kappa} - \frac{1}{a^3}\right) \right\} \right] \quad (\text{A.4}) \end{aligned}$$

The Kramers-Moyal expansion is expressed as

$$\begin{aligned}
 & \exp\left[\pm \left\{ F(n_\nu, n_\kappa) - F\left(n_\nu + \frac{1}{a^3}, n_\kappa - \frac{1}{a^3}\right) \right\} / 2k_B T\right] \\
 &= 1 \mp \frac{1}{a^3} \frac{1}{2k_B T} \left(\frac{\partial}{\partial n_\nu} - \frac{\partial}{\partial n_\kappa} \right) F \mp \frac{1}{2a^6} \frac{1}{2k_B T} \left(\frac{\partial}{\partial n_\nu} - \frac{\partial}{\partial n_\kappa} \right)^2 F \\
 &+ \frac{1}{2a^6} \left(\frac{\partial F}{\partial n_\nu} - \frac{\partial F}{\partial n_\kappa} \right)^2 \left(\frac{1}{2k_B T} \right)^2 + \dots
 \end{aligned}$$

(A.5)

With eq. (2.A.5), eq. (2.A.4) is modified to

$$\begin{aligned}
 \frac{\partial P(\{n_\mu\})}{\partial t} &= - \frac{1}{2\tau a^2} \sum_{\langle \nu \kappa \rangle} \frac{1}{k_B T} \left[\frac{\partial}{\partial n_\nu} \left\{ (2n_\nu(1-n_\nu) + (n_\nu - n_\kappa)(2n_\nu - 1)) \right. \right. \\
 &\quad \left. \left. \times \left\{ \left(\frac{\partial}{\partial n_\nu} - \frac{\partial}{\partial n_\kappa} \right) F \right\} P \right\} \right. \\
 &\quad \left. + \frac{1}{2} k_B T \left\{ \left(\frac{\partial}{\partial n_\nu} - \frac{\partial}{\partial n_\kappa} \right) \sqrt{n_\kappa(1-n_\kappa)n_\nu(1-n_\nu)} \left(\frac{\partial}{\partial n_\nu} - \frac{\partial}{\partial n_\kappa} \right) P \right\} \right] \quad (A.6)
 \end{aligned}$$

where only the terms up to the first order of $(n_\nu - n_\kappa)$ are retained.

Thus we obtain a Fokker-Planck equation. If we assume that the fluctuation is small, we can consider only the first term in the right hand side of eq. (A.6). Thus we obtain

$$\begin{aligned}
 \frac{\partial P(\{n_\mu\})}{\partial t} &= - \frac{1}{2\tau a^2 k_B T} \sum_{\langle \nu \kappa \rangle} \frac{\partial}{\partial n_\nu} \left\{ (2n_\nu(1-n_\nu) + (n_\nu - n_\kappa)(2n_\nu - 1)) \right. \\
 &\quad \left. \times \left\{ \left(\frac{\partial}{\partial n_\nu} - \frac{\partial}{\partial n_\kappa} \right) F \right\} P \right\} \quad (A.7)
 \end{aligned}$$

If we multiply eq. (2.A.8) by n_μ , and integrate by parts, we obtain

$$n_\mu \frac{\partial P}{\partial t} = \frac{1}{2\tau a^2 k_B T} \sum_\delta \left\{ 2n_\mu(1-n_\mu) + (n_\mu - n_{\mu+\delta})(2n_\mu - 1) \right\} \\ \times \left\{ \left(\frac{\partial}{\partial n_\mu} - \frac{\partial}{\partial n_{\mu+\delta}} \right) F \right\} P, \quad (A.8)$$

where δ denotes a co-ordinate of the nearest neighbour cell.

With the gradient expansions of $(n_\mu - n_{\mu+\delta})$ and $\left(\frac{\partial}{\partial n_\mu} - \frac{\partial}{\partial n_{\mu+\delta}} \right)$, eq. (A.8) is transformed into

$$n_\mu \frac{\partial P}{\partial t} = - \frac{1}{\tau k_B T} \nabla \left\{ (n_\mu(1-n_\mu)) \cdot \nabla \frac{\partial F}{\partial n_\mu} \right\} P. \quad (A.9)$$

An ensemble average of a variable X is denoted by $\langle X \rangle =$

$\sum_{\{n_i\}} P(\{n_i\}) X(\{n_i\})$. If we neglect the fluctuation $\langle X^2 \rangle - \langle X \rangle^2$, and use the conservation law

$$\frac{\partial}{\partial t} \left[\sum_{\{n_i\}} n_\mu P \right] = 0, \quad (A.10)$$

we obtain from eq. (A.9)

$$\frac{\partial \langle n_\mu \rangle}{\partial t} = \frac{1}{\tau k_B T} \nabla \left\{ \langle n_\mu \rangle (1 - \langle n_\mu \rangle) \cdot \nabla \left\langle \frac{\partial F}{\partial n_\mu} \right\rangle \right\}. \quad (A.11)$$

This is equivalent to eq. 3.1).

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Figure caption

Fig. 1. The growth of modes. The curve represents the growth of the concentration inhomogeneity as a function of the largest growing mode. Time is a parameter.

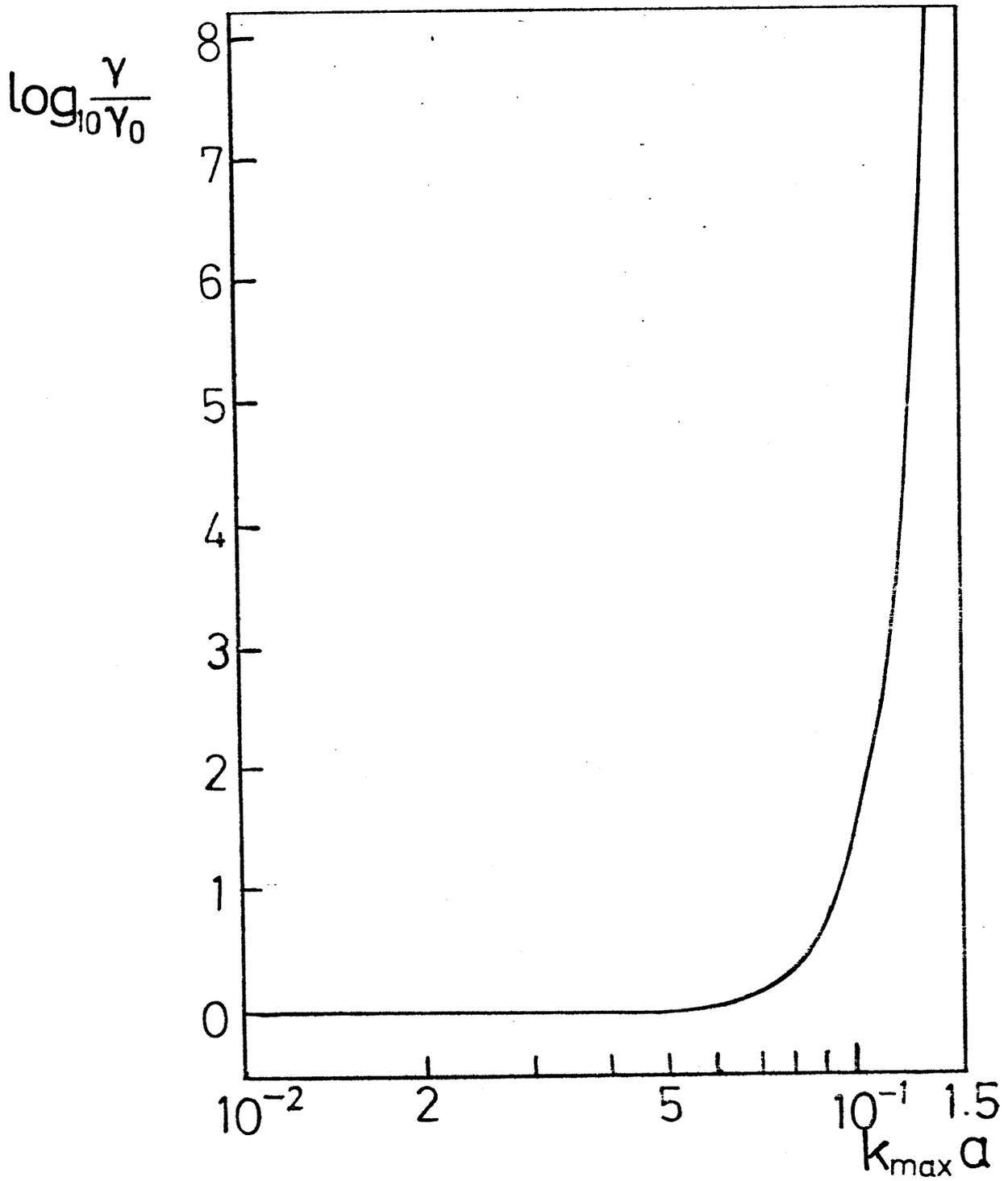


Fig. 1

Table I

The binding energy of a divacancy

Metal	Binding energy	Remark	Reference
Au	<0.16 eV	Experiment	a)
Ag	0.24 eV	Experiment	a)
Cu	0.12 eV	Experiment	a)
Ni	0.28 eV	Experiment	b)
Cu	0.29 eV	Theory	13)
Ag	0.33 eV	Theory	13)
Au	0.28 eV	Theory	13)
Ni	0.4~0.5 eV	Theory	13)

a) H.Mehrer and A.Seeger: Phys. Stat. Solidi 39(1970)647.

b) A.seeger, G.schottky and D.Schumacher: Phys. Stat. Solidi 11 (1965)363.

Table II^{*)}

Metal	Formation Energy (eV)	Metal	Formation Energy (eV)
Au	0.87	Nb	2.0
Ag	0.99	Mo	2.3
Cu	1.03	W	3.3
Al	0.73		

*) Ref.16 p.27