6 Dislocations in Other Crystal Structures

6.1 Introduction
The face-centred cubic metals have been treated in a separate chapter because although many of the dislocation reactions and properties presented have counterparts in other structures, they are more readily described in the face-centred cubic system. Also, these metals have been more extensively studied than other solids. In general, reducing the crystal symmetry, changing the nature of the interatomic bonding and increasing the number of atom species in the lattice make dislocation behaviour more complex. Nevertheless, many of the features of the preceding chapter carry over to other structures, as will be seen in the following. The two other major metallic structures are discussed first, and then some important compounds and non-metallic cases are considered.

6.2 Dislocations in Hexagonal Close-packed Metals

Burgers Vectors and Stacking Faults
Some of the important metals of this structure are given in Table 6.1. As explained in section 1.2, the (0001) basal plane is close-packed and the close-packed directions are \(\{11\bar{2}0\}\). The shortest lattice vectors are \(\frac{1}{3}(11\bar{2}0)\), the unit cell generation vectors \(a\) in the basal plane. It may be anticipated, therefore, that dislocation glide will occur in the basal plane with Burgers vector \(b = \frac{1}{3}(11\bar{2}0)\); this slip system is frequently observed. None of the metals is ideally close-packed, which would require the lattice parameter ratio \(c/a = (8/3)^{1/2} = 1.633\), although magnesium and cobalt have a \(c/a\) ratio close to ideal. This suggests that directionality occurs in interatom bonding. In support of this, it is found that some metals slip most easily with \(b = \frac{1}{3}(11\bar{2}0)\) on the first-order prism planes \\{10\bar{1}0\} (see Table 6.1).

Burgers vectors for the structure may be described in a similar fashion to the Thompson tetrahedron for face-centred-cubic metals by using the bi-pyramid shown in Fig. 6.1. The important dislocations and their Burgers vectors are as follows.

| Table 6.1 Properties of some hexagonal close-packed metals at 300 K |
|-----------------|---|---|---|---|---|---|---|
| Metal           | Be | Ti | Zr | Mg | Co | Zn | Cd |
| c/a ratio       | 1.568 | 1.587 | 1.593 | 1.623 | 1.628 | 1.856 | 1.886 |
| Preferred slip  | basal | prism | prism | basal | basal | basal | basal |
| Plane for \(b = a\) | (0001) | \{10\bar{1}0\} | \{10\bar{1}0\} | (0001) | (0001) | (0001) | (0001) |
Dislocations in other crystal structures

(a) Perfect dislocations with one of six Burgers vectors in the basal plane along the sides of the triangular base $ABC$ of the pyramid, represented by $AB$, $BC$, $CA$, $BA$, $CB$ and $AC$.

(b) Perfect dislocations with one of two Burgers vectors perpendicular to the basal plane, represented by the vectors $ST$ and $TS$.

(c) Perfect dislocations with one of twelve Burgers vectors represented by symbols such as $SA/TB$, which means either the sum of the vectors $ST$ and $AB$ or, geometrically, a vector equal to twice the join of the midpoints of $SA$ and $TB$.

(d) Imperfect basal dislocations of the Shockley partial type with Burgers vectors $A\sigma$, $B\sigma$, $C\sigma$, $\sigma A$, $\sigma B$ and $\sigma C$.

(e) Imperfect dislocations with Burgers vectors perpendicular to the basal plane, namely, $\sigma S$, $\sigma T$, $S\sigma$ and $T\sigma$.

(f) Imperfect dislocations which are a combination of the latter two types given by $AS$, $BS$, etc. Although these vectors represent a displacement from one atomic site to another the associated dislocations are not perfect. This is because the sites do not have identical surroundings and the vectors are not translations of the lattice.

The Miller-Bravais indices and length $b$ of these Burgers vectors $b$ are given in Table 6.2. The value of $b^2$ for ideal close-packing ($c^2 = \frac{3}{2}a^2$) is also given: appropriate adjustments are required when dealing with non-ideal packing.

A number of different stacking faults are associated with the partial dislocations listed in Table 6.2. According to the hard-sphere model of atoms, three basal-plane faults exist which do not affect nearest-neighbour arrangements of the perfect stacking sequence $ABABAB\ldots$
Table 6.2 Dislocations in hexagonal close-packed structures

<table>
<thead>
<tr>
<th>Type</th>
<th>AB</th>
<th>TS</th>
<th>SA/TB</th>
<th>Aσ</th>
<th>σS</th>
<th>AS</th>
</tr>
</thead>
<tbody>
<tr>
<td>b</td>
<td>$\frac{1}{3}(1\bar{1}20)$</td>
<td>[0001]</td>
<td>$\frac{1}{3}(1\bar{1}23)$</td>
<td>$\frac{1}{3}(1\bar{1}00)$</td>
<td>$\frac{1}{2}$[0001]</td>
<td>$\frac{1}{6}(2203)$</td>
</tr>
<tr>
<td>b</td>
<td>a</td>
<td>c</td>
<td>$(c^2 + a^2)^{1/2}$</td>
<td>$a/\sqrt{3}$</td>
<td>c/2</td>
<td>$(a^2 + c^2)^{1/2}$</td>
</tr>
<tr>
<td>$b^2$</td>
<td>$a^2$</td>
<td>$\frac{8}{3}a^2$</td>
<td>$\frac{11}{3}a^2$</td>
<td>$\frac{1}{3}a^2$</td>
<td>$\frac{2}{3}a^2$</td>
<td>$a^2$</td>
</tr>
</tbody>
</table>

Two are intrinsic and conventionally called $I_1$ and $I_2$. Fault $I_1$ is formed by removal of a basal layer, which produces a very high energy fault, followed by slip of $\frac{1}{3}(1\bar{1}0\bar{1})$ of the crystal above this fault to reduce the energy:

$$ABABABABA\ldots \rightarrow ABABBABA\ldots \rightarrow ABABCBCB\ldots (I_1)$$

Fault $I_2$ results from slip of $\frac{1}{3}(10\bar{1}0)$ in a perfect crystal:

$$ABABABAB\ldots \rightarrow ABABCACA\ldots (I_2)$$

The extrinsic fault $(E)$ is produced by inserting an extra plane:

$$ABABABAB\ldots \rightarrow ABABCABAB\ldots (E)$$

These faults introduce into the crystal a thin layer of face-centred cubic stacking $(ABC)$ and so have a characteristic stacking-fault energy $\gamma$. The main contribution to $\gamma$ arises from changes in the second-neighbour sequences of the planes. There is one change for $I_1$, two for $I_2$ and three for $E$, and so to a first approximation $\gamma_E \approx \frac{1}{3}\gamma_{I_2} \approx 3\gamma_{I_1}$. Experimental estimates of $\gamma$ for the basal-slip metals are generally higher than those quoted earlier for the face-centred cubic metals. Stable faults for the basal plane of the prism-slip metals such as titanium and zirconium either have very high energy or may not even exist.

There is no direct experimental evidence of stacking faults on other planes, which are not close-packed; but the incidence of non-basal glide and results of simulations using computer models suggest that stable faults may exist on prism and pyramidal planes, although possibly with high energy.

**Basal and Non-basal Slip**

Slip by the system $\frac{1}{3}(1\bar{1}20)(0\bar{0}01)$ in metals such as beryllium, magnesium, cadmium and zinc is similar to $\frac{1}{3}(11\bar{1}0)\{11\bar{1}\}$ slip in the face-centred cubic metals, in that the critical resolved shear stress is low ($\leq 1$MN m$^{-2}$) and the perfect dislocation dissociates into two Shockley partials bounding a ribbon of stacking fault. The Burgers vector reaction is

$$AB \rightarrow A\sigma + \sigma B$$

(6.4)
Dislocations in other crystal structures

Figure 6.2 Dissociation of a perfect dislocation with Burgers vector $\mathbf{AB}$ (Fig. 6.1) into two Shockley partial dislocations separated by a stacking fault $I_2$. Double arrows indicate the two errors in the two-fold stacking sequence of the basal planes.

e.g.

$$\frac{1}{3}[1120] \rightarrow \frac{1}{3}[10\overline{1}0] + \frac{1}{3}[01\overline{1}0]$$

$$b^2: \frac{a^2}{3} \quad \frac{a^2}{3} \quad \frac{a^2}{3}$$

The geometry is the same as that of the face-centred cubic case, in that the partial vectors lie at $\pm 30^\circ$ to the perfect vector and the fractional reduction in dislocation energy given by $b^2$ is $1/3$. The spacing of the partials may be calculated from equations (5.4) and (5.5) by replacing $b^2$ by $a^2/3$. The fault involved is the intrinsic fault $I_2$, as shown schematically in Fig. 6.2.

Basal slip also occurs in the metals for which prism slip is most favoured, but the critical resolved shear stress is high ($\approx 100$ MN m$^{-2}$). It is probable, therefore, that interatomic forces prevent formation of a stable basal fault and that dissociation (6.4) does not take place.

Slip on non-basal systems occurs most commonly with $\mathbf{b} = \frac{1}{4}[12\overline{1}0]$. It is seen from Fig. 6.3 that dislocations with this slip vector can glide on the first-order prism and pyramidal planes in addition to the basal plane. However, in metals for which basal slip is preferred, the resolved shear stress required for non-basal slip is one to two orders of magnitude higher. This is explained by the absence of low-energy stacking faults on other planes and the subsequent difficulty of dislocation glide. Furthermore, a dislocation extended in the basal plane by reaction (6.4) can only cross-slip onto other planes by first constricting, as explained in section 5.3. This will be hindered by the absence of other faults except at high values of $\tau_B$, stress and temperature.

Even in metals for which it is the favoured slip system, the critical resolved shear stress for $\frac{1}{4}[1120]\{1100\}$ prism slip is high ($\gtrsim 10$ MN m$^{-2}$). An obvious consideration for basal and prism slip is to compare the
inter-planar spacings of the two sets of slip planes, for in simple terms the lattice resistance to glide tends to be smaller for planes with wide spacings (section 10.3). The spacings of the basal and corrugated prism planes are $c/2$ and $\sqrt{3}a/2$ respectively, and thus the prism-plane spacing is greater when $c/a < \sqrt{3}$. This approach fails to explain the preference for basal slip in several metals (Table 6.1) and points to the importance of stacking faults in determination of slip systems. Two dissociations have been proposed to account for prism slip in titanium and zirconium, namely

\begin{equation}
\frac{1}{3} \langle 11\bar{2}0 \rangle \rightarrow \frac{1}{18} \langle 42\bar{6}3 \rangle + \frac{1}{18} \langle 24\bar{6}3 \rangle
\end{equation}

(6.5)

\begin{equation}
\frac{1}{3} \langle 11\bar{2}0 \rangle \rightarrow \frac{1}{9} \langle 11\bar{2}0 \rangle + \frac{2}{9} \langle 11\bar{2}0 \rangle
\end{equation}

(6.6)

The first produces a fault which is stable in a lattice consisting of hard spheres. The second results in prism planes adjacent to the fault ribbon adopting the stacking of \{112\} planes in body-centred cubic metals, the relevance being that this is the stable crystal structure for titanium and zirconium at high temperature. There is no evidence that these faults exist in real metals, however, and lattice imaging using transmission electron microscopy shows that dislocation cores do not extend widely on the prism plane. This is supported by atomic-scale computer modelling, which also indicates that the core with a narrow width on the prism plane is favoured energetically over the dissociated core on the basal plane (equation 6.4) because of the high energy of the $I_2$ basal fault in these metals.

In polycrystalline metals, basal and prism slip do not supply sufficient slip modes to satisfy von Mises’ criterion that every grain should be able to plastically deform generally to meet the shape changes imposed by its neighbours (section 10.9). This requires five independent slip systems, whereas basal and prism slip provide only two each. Consequently, twinning and occasionally other slip systems play an important role in the plasticity of these metals.

A variety of twin modes with different habit plane and shear direction occur, depending on the metal, temperature and nature of loading. A deformation twin grows or shrinks under a resolved shear stress by the motion of steps along its boundary. Such steps have dislocation character and are known as twinning dislocations. They are described in more detail in section 9.7.

Slip with Burgers vector $\frac{1}{3} \langle 11\bar{2}3 \rangle$ has been widely reported, although only under conditions of high stress and orientations in which more favoured slip vectors cannot operate. The glide planes are $\{10\bar{1}1\}$ and $\{11\bar{2}2\}$. The magnitude of the Burgers vector is large and the planes are atomically rough, which explains the high glide stress. It has been proposed on the basis of the hard-sphere lattice model that the dislocation on $\{11\bar{2}2\}$ may dissociate into four partials with approximately equal
Burgers vectors. One has a core spread over three successive \{11\over22\} planes, within which the atomic displacements are different. These three adjacent faults are bounded by the three remaining partials. Significantly, the first dislocation is a zonal twinning dislocation, a multi-layer step in a coherent twin boundary which generates by its movement a twin. Thus, there may be a close relationship between \{11\over23\}{11\over22} slip and twinning, although not necessarily of the hard-sphere lattice form.

**Vacancy and Interstitial Loops**

As in the face-centred cubic metals, vacancies and interstitials in excess of the equilibrium concentration can precipitate as platelets to form dislocation loops. The situation is more complicated, however, because although the basal plane is close-packed, the relative density of atoms in the different crystallographic planes varies with \(c/a\) ratio. Also, stacking faults occur in some metals and not others, as discussed in the preceding section. The simplest geometries are considered here, starting with basal-plane loops.

Condensation of vacancies in a single basal plane (Fig. 6.4(a)) results in two similar atomic layers coming into contact (Fig. 6.4(b)). This unstable situation of high energy is avoided in one of two ways. In one, the stacking of one layer adjacent to the fault is changed; for example \(B\) to \(C\) as in Fig. 6.4(c). This is equivalent to the glide below the layer of one Shockley partial with Burgers vector \(\frac{1}{2}(10\overline{1}0)\) followed by glide above of a second Shockley of opposite sign. The Burgers vector reaction is

\[
\sigma S + \sigma A + \Lambda \sigma \rightarrow \sigma S
\]

![Diagram](image)

**Figure 6.4** Formation of prismatic dislocation loops as a result of the precipitation of one layer of vacancies. (a) Disc-shaped cavity. (b) Collapse of the disc bringing two \(B\) layers together. (c) Formation of a high-energy stacking fault \(E\). (d) Formation of a low-energy stacking fault \(I\). The actual sequence of planes formed will depend on the plane in which the vacancies form, i.e. \(A\) or \(B\). It is possible to have loops of the same form with different Burgers vectors lying on adjacent planes. (After Berghezan, Fourdeux and Amelinckx, *Acta Metall.* 9, 464, 1961.)
The resultant sessile Frank partial with \( b^2 = c^2/4 \approx 2a^2/3 \) surrounds the extrinsic (E) stacking fault described earlier in this section. In the alternative mechanism, a single Shockley partial sweeps over the vacancy platelet, displacing the atoms above by \( \frac{1}{3} \langle 1010 \rangle \) relative to those below. The Burgers vector reaction is
\[
A \sigma + \sigma S \rightarrow AS
\]
\[\text{e.g.}\]
\[
\frac{1}{2} [0001] + \frac{1}{3} [1\bar{1}00] + \frac{1}{3} [\bar{1}100] \rightarrow \frac{1}{2} [0001]
\] (6.7)

This sessile Frank partial with \( b^2 \approx a^2 \) surrounds the type \( I_1 \) intrinsic fault (Fig. 6.4(d)). The \( E \)-type loop of reaction (6.7) can transform to the \( I_1 \) form by reaction (6.8), and since \( \gamma_E \) is expected to be approximately three times \( \gamma_{I_1} \), the \( \frac{1}{6} \langle 2203 \rangle \) loops may be expected to dominate. The dislocation energy is proportional to \( b^2 \), however, and the total energy change accompanying reaction (6.8) is dependent on the \( \gamma \) values and loop size, in a similar manner to the unfaulting of loops in face-centred cubic metals (section 5.5). There is therefore a critical loop size for the

![Figure 6.5](image_url)

**Figure 6.5** (a) Precipitation of a layer of interstitials. (b) Prismatic loop resulting from the layer of interstitials: the loop contains a high-energy stacking fault \( E \). (c) Prismatic loop containing low-energy stacking fault \( I_1 \). (After Berghezan, Fourdeux and Amelinckx, *Acta Metall.* 9, 464, 1961.)
reaction, which may be influenced by factors such as stress, temperature and impurity content. Experimentally, both forms are observed in quenched or irradiated metals.

Precipitation of a basal layer of interstitials as in Figs 6.5(a) and (b) produces an $E$-type fault surrounded by a Frank partial loop with Burgers vector $\frac{1}{2}[0001]$. Again, provided the loop is large enough, this can transform to an $I_1$ loop by the nucleation and sweep of a Shockley partial according to reaction (6.8). The stacking sequence is shown in Fig. 6.5(c). Interstitial loops with Burgers vectors $\frac{1}{2}[0001]$ and $\frac{1}{6}(20\bar{2}3)$ have been seen in irradiated magnesium, cadmium and zinc. In the latter two metals, perfect loops with Burgers vector [0001] are also observed. They result from a double layer of interstitials, and as they grow during irradiation, some loops transform into two concentric $\frac{1}{2}[0001]$ loops, as shown in Fig. 6.6.

The atomic density of the basal planes is only greater than that of the corrugated $\{10\bar{1}0\}$ prism planes when $c/a > \sqrt{3}$, suggesting that the existence of basal stacking faults aids the stability of the basal vacancy and interstitial loops in magnesium. These faults are not believed to occur in titanium and zirconium, however, and large basal-plane loops are not expected. This is confirmed experimentally by transmission-electron microscopy of irradiated specimens, which reveals that basa-

Figure 6.6  Transmission electron micrographs of interstitial loops with $b = [0001]$ in zinc. (a) Below the threshold voltage at which electrons displace zinc atoms and (b) above. Loops $\beta$ split into concentric loops $\gamma$ as they grow. (From Eyre, Loretto and Smallman (1976), *Proc. Vacancies '76*, p. 63, The Metals Society, London.)
plane loops with a $c$ component in their Burgers vector occur only rarely, for example at high levels of radiation damage. Generally, both vacancy and interstitial loops with the perfect Burgers vector $\frac{1}{2} \langle 1120 \rangle$ are seen. They lie on planes of the $\{0001\}$ zone at angles up to $30^\circ$ from the pure-edge $\{1120\}$ orientation. The $\{1120\}$ planes are neither widely-spaced nor densely-packed, and it is possible that the point defects precipitate initially as single-layer loops on the $\{1010\}$ prism planes. The resulting stacking fault probably has a high energy and is removed by shear when the loops are small by the Burgers vector reaction

$$ \frac{1}{2} \langle 10\overline{1}0 \rangle + \frac{1}{6} \langle \overline{1}2\overline{1}0 \rangle \rightarrow \frac{1}{3} \langle 11\overline{2}0 \rangle $$

(6.9)

The glissile loops thus produced can adopt the variety of orientations observed in practice. The small loops shown in Fig. 2.15 are examples of vacancy loops with $b = \frac{1}{2} \langle 1010 \rangle$ or $\frac{1}{3} \langle 1120 \rangle$ in ruthenium irradiated with heavy ions.

### 6.3 Dislocations in Body-centred Cubic Metals

In body-centred cubic metals (e.g. iron, molybdenum, tantalum, vanadium, chromium, tungsten, niobium, sodium and potassium) slip occurs in close-packed $\langle 111 \rangle$ directions. The shortest lattice vector, i.e. the Burgers vector of the perfect slip dislocation, is of the type $\frac{1}{2} \langle 111 \rangle$. The crystallographic slip planes are $\{110\}$, $\{112\}$ and $\{123\}$. Each of these planes contains $\langle 111 \rangle$ slip directions and it is particularly significant that three $\{110\}$, three $\{112\}$ and six $\{123\}$ planes intersect along the same $\langle 111 \rangle$ direction. Thus, if cross slip is easy it is possible for screw dislocations to move in a haphazard way on different $\{110\}$ planes or combinations of $\{110\}$ and $\{112\}$ planes, etc., favoured by the applied stress. For this reason slip lines are often wavy and ill-defined. It has been found that the apparent slip plane varies with composition, crystal orientation, temperature and strain rate. Thus, when pure iron is deformed at room temperature the slip plane appears to be close to the maximum resolved shear stress plane irrespective of the orientation, whereas when it is deformed at low temperatures, or alloyed with silicon, slip tends to be restricted to a specific $\{110\}$ plane.

An interesting feature of yielding is the asymmetry of slip. It is found, for example, that the slip plane of a single crystal deformed in uniaxial compression may be different from the slip plane which operates in tension for the same crystal orientation. In other words, the shear stress to move a dislocation lying in a slip plane in one direction is not the same as the shear stress required to move it in the opposite direction in the same plane. Slip is easier when the applied stress is such that a dislocation would move in the twinning sense (see below) on $\{112\}$ planes rather than the anti-twinning sense, even when the actual slip plane is not $\{112\}$. Electron microscopy of metals deformed at low temperature reveals long screw dislocations, implying that non-screw dislocations are more mobile and that screw dislocations dictate the slip characteristics.
It is a feature of the body-centred cubic metals that stacking faults have not been observed experimentally, and the ease of cross slip suggests that faults are at best of very high energy. Consequently, simple elastic calculations of dislocation dissociation are inadequate and much of our knowledge about core structure and behaviour has been obtained from computer models of crystals, in which atoms are allowed to interact by suitable interatomic potentials (see section 2.7). By computing the energy of such crystals when stacking faults are deliberately created on the low-index planes, it has been confirmed that stable faults are unlikely to exist in this structure.

Computer simulations of the screw dislocation with Burgers vector \( \frac{1}{2}(111) \) show that the core has a non-planar character. An example of two equivalent atomic configurations for a \( \frac{1}{2}(111) \) screw are shown in Figs 6.7(a) and (b). The small circles represent atom positions projected on the (111) plane: the dislocation is perpendicular to the plane at the centre of each diagram. The orientation of the traces of the \{110\} and \{112\} planes are shown in Fig. 6.7(c). The only atomic displacements which are not negligible are parallel to the dislocation line [111], and the atom projections are the same as for a perfect crystal. In order to represent the [111] disregistry of atoms in the core, arrows are drawn between pairs of neighbouring atoms on the projections. The length of an arrow is proportional to the difference of the [111] displacements of the two atoms, and scaled such that for a displacement difference of \( b/3 \), the arrow just runs from one atom to the other. When the difference falls between \( b/2 \) and \( b \), it is reduced by \( b \). For the isotropic elastic solution (equation (4.11)), the length of the arrows would decrease in inverse proportion to the distance from the core centre and would exhibit complete radial symmetry. In the atomic model, the displacements are concentrated on the three intersecting \{110\} planes, each of which contains an unstable fault produced by a \( \frac{1}{2}(111) \) displacement. Although the \( \frac{1}{2}(111) \) dislocation spreads into three \( \frac{1}{6}(111) \) cores, these fractional dislocations do not bound stable stacking faults, unlike Shockley partial dislocations. Close examination of the displacements reveals that the fractional cores also spread asymmetrically, but on three \{112\} planes in the twinning sense. This accounts for the slip asymmetry referred to above.

By applying stress to the model crystals, it has been found that the core structure changes before slip occurs. For example, under an increasing shear stress on (101) tending to move the screw dislocation in Fig. 6.7(a) to the left, the fractional dislocation on (101) extends the core to the left and the two others constrict towards the core centre, as shown in Fig. 6.8(a). As the stress is increased (Fig. 6.8(b)), the fractional dislocation on (011) disappears to be replaced by another on (110) before glide of the whole core occurs on (101). Movement of the core to the right under stress takes a different form. Although the detailed core changes are dependent on the interatomic potentials, it is found that under pure shear stress slip occurs on the \{110\} planes with the
asymmetry described above. Furthermore, computer simulation has shown that the screw core responds differently to stresses with different non-shear components, in good agreement with the effects of compression and tension found in experiment.

Similar studies of non-screw dislocations, on the other hand, show they have cores which are planar in form on either \{110\} or \{112\} but do not contain stable stacking faults. Like their face-centred cubic counterparts, they are not sensitive to the application of non-shear stresses, and they glide at much lower shear stresses than the screw dislocation.

Another set of perfect dislocations in the body-centred cubic metals are those with Burgers vector \(\frac{1}{2}[111]\). They are occasionally observed in
dislocations in other crystal structures

Figure 6.8  Core structures of the dislocation of Fig. 6.7(a) under applied shear stress of (a) 0.0115G and (b) 0.0265G, where G is the shear modulus. The whole dislocation moves at a stress of 0.0275 G. (After Duesbery, Vitek and Bowen, Proc. Roy. Soc. A332, 85, 1973.)

dislocation networks and are believed to occur from the reaction of two perfect \( \frac{1}{2} \langle 111 \rangle \) dislocations:

\[
e.g. \quad \frac{1}{2} [111] + \frac{1}{2} [1\overline{1}] \rightarrow [100]
\]

(6.10)

Computer modelling of the edge dislocation has shown that large tensile forces in the core region just below the extra half-plane lead to severing of atomic bonds and the formation of a microcrack there. This dislocation is therefore unlikely to take part in plastic deformation, but since the main cleavage planes are \( \{100\} \), it may play a role in crack nucleation.

Deformation twinning is observed in all the body-centred cubic transition metals when they are deformed at low temperature and/or high strain rate. As noted with respect to the hexagonal close-packed metals in section 6.2, and described in more detail in section 9.7, the atomic displacements that allow growth of a twin occur by the glide of steps (twinning dislocations) over the twin boundary plane. For the body-centred cubic metals twinning occurs on \( \{112\}\{111\} \) systems. In section 1.2 it was shown that the stacking sequence of \( \{112\} \) planes in the body-centred cubic structure is \( ABCDEFA... \). The homogeneous shear required to produce a twin is \( 1/\sqrt{2} \) in a \( \langle 111 \rangle \) direction on a \( \{112\} \) plane. This shear can be produced by a displacement of \( \frac{1}{6} \langle 111 \rangle \) on every successive \( \{112\} \) plane i.e. the Burgers vector of a twinning dislocation has the form \( \frac{1}{6} \langle 111 \rangle \) and the step height is \( a/\sqrt{6} \). It is seen from Fig. 1.6(b) that if all the atoms in, say, an \( E \) layer and above are translated by \( \frac{1}{6} \langle 111 \rangle \), then \( E \) shifts to \( C \), \( F \) shifts to \( D \), etc., and the new sequence is \( ABCDCDE... \). A second translation on the adjacent plane displaces \( D \) to \( B \), \( E \) to \( C \), etc., resulting in \( ABCDCBC... \). Repetition of this translation...
on successive planes gives \( ABCDCBA \ldots \), which is the stacking of a twinned crystal. (This process is described in more detail in section 9.7.)

From the atom positions in Fig. 1.6(b), it is clear that the single translation \( \frac{1}{6}[\overline{1}11] \) which displaces \( E \) to \( C \), \( F \) to \( D \), etc., produces a different result from the translation \( \frac{1}{6}[111] \) of opposite sense. The latter moves \( E \) to \( D \) and produces an untwinned structure of high energy. There is thus an asymmetry with respect to \( \frac{1}{6}(111) \) translations in the twinning and anti-twinning sense on \( \{112\} \) planes.

Dislocation loops formed by interstitials and vacancies are an important product of radiation damage. There are no close-packed planes in the body-centred cubic structure, and it has been suggested that the loops nucleate with Burgers vector \( \frac{1}{2}(110) \) on the \( \{110\} \) planes, which are the most densely packed. In the absence of stable stacking faults on these planes, the partial dislocation loops would shear at an early stage of growth to become perfect by one of two reactions:

\[
\begin{align*}
\frac{1}{2}(110) + \frac{1}{2}(001) & \rightarrow \frac{1}{2}(111) \\
\frac{1}{2}(110) + \frac{1}{2}(1\overline{1}0) & \rightarrow (100)
\end{align*}
\]

The resultant dislocation in reaction (6.11) has the lower energy (i.e. \( b^2 \)) and loops with this Burgers vector have been observed in many metals. Computer simulation (section 2.7) of damage production in displacement cascades shows that small clusters of self-interstitial atoms nucleate directly as loops with Burgers vector \( \frac{1}{2}(111) \). A cluster consists of \( (111) \) crowdion interstitials, i.e. each interstitial is an extra atom inserted in a close-packed \( (111) \) atomic row, as illustrated by the computer-generated image of a cluster of 19 interstitials in Fig. 6.9. The distortion is focused along the crowdion axis and simulation predicts that these small loops move easily along their glide prism. The one in Fig. 6.9 has moved by about \( 7a \).

In \( \alpha \)-iron and its alloys, however, a high proportion of loops are of \( (100) \) type: this somewhat surprising effect is as yet unexplained.

### 6.4 Dislocations in Ionic Crystals

An important feature of dislocations in ionic solids is that electrical charge effects can be associated with them. For example, compressive deformation increases electrical conductivity and electric current is produced during plastic deformation by the motion of charged dislocations. Ionic crystals contain atoms of elements from different sides of the periodic table which transfer electrons from one species to the other, producing sets of cations and anions. One of the simplest forms is the rocksalt (\( NaCl \)) structure shown in Fig. 1.12, in which each anion is surrounded by six cations and vice versa. It has a face-centred cubic Bravais lattice with an anion–cation pair for each lattice point, one ion at \( 0, 0, 0 \) and the other at \( \frac{1}{2}, 0, 0 \). MgO, LiF and AgCl also have this
structure. It has been widely studied and its description provides a basis for more complicated systems.

The shortest lattice vector is $\frac{1}{2}(110)$, and this is the Burgers vector of the dislocations responsible for slip. The principal slip planes are {110}. Slip steps are also observed on {100} and (occasionally) {111} and {112} planes after high stresses, particularly at high temperatures and in crystals of high polarizability, where the ionic nature of bonding decreases. Cross slip of screw dislocations can occur only by glide on
planes other than \{110\}, for only one \(\langle 110 \rangle\) direction lies in a given \{110\) plane.

Figure 6.10 shows a pure edge dislocation with a \(\frac{1}{2}\langle 110 \rangle\) Burgers vector and \(\langle 110 \rangle\) slip plane emerging on the \(001\) surface. The extra half-plane actually consists of two supplementary half-planes, as shown. The ions in the planes below the surface alternate between those shown in Figs 6.10(a) and (b). The figure serves to illustrate that there is an effective charge associated with the point of emergence of the dislocation on the \(001\) surface. Intuitively, if it is \(-q\) in (a) it must be \(+q\) in (b). It is readily shown that \(q = e/4\), where \(e\) is the electronic charge, as follows.

In any cube of NaCl in which the corner ions are of the same type, as in Fig. 1.12, the ions of that sign exceed in number those of opposite sign by one. This excess charge of \(\pm e\) may be considered as an effective charge \(\pm e/8\) associated with each of the eight corners. In a cube with equal numbers of anions and cations, the positive and negative corners neutralise each other. Consider the dislocation of Fig. 6.10(a) to be in a block of crystal \(ABCD\) bounded by \{100\} faces. The effective charge of the four corners \(A, B, C, D\) is \((+e/8 - 3e/8) = -e/4\), and so the net effective charge at the \(001\) surface is \((-e/4 - q)\). Remove a single layer of ions to expose the new face \(A'B'C'D'\) (Fig. 6.10(b)). The net effective charge is now \((-e/8 + 3e/8 + q)\), which has been achieved by removing sixteen anions and fifteen cations, i.e. \(-e\). Since the initial charge must equal the final charge plus the charge removed,

\[
\frac{-e}{4} - q = \frac{e}{4} + q - e
\]

and \(q = e/4\). The same result holds for emergence on \{110\) planes.
When the edge dislocation glides on its (110) slip plane, there is no displacement (and hence transport of charge) along the line, and the effective charge of the emergent point does not change sign. For the same reason, *kinks* in edge dislocations bear no effective charge. If the dislocation climbs by, say, removal of the anion labelled \( n \) at the bottom of the extra half-plane in Fig. 6.10(a), the configuration changes to the mirror image of that in Fig. 6.10(b), and so the effective charge at the emergent point changes sign. It follows that *jogs* carry effective charge, as demonstrated by the illustration in Fig. 6.11 of two *elementary* (or *unit*) jogs of one atom height. The bottom row of ions has an excess charge \(-e\), which is effectively carried by the two jogs. Thus, depending on the sign of the end ion of the incomplete row, an elementary jog has a charge \( \pm e/2 \). It cannot be neutralized by point defects of integer charge. Charged jogs attract or repel each other electrostatically, and only jogs of an even number times the elementary height are neutral. (Note that in divalent crystals such as MgO, the effective charges are twice those discussed here.)

The formation energies of anion and cation vacancies are in general different, and this results in a higher probability of a jog being adjacent to a vacant site of the lower energy. Dislocations thus have an effective charge per unit length in thermal equilibrium, although this is neutralised in the crystal overall by an excess concentration of vacancies of opposite sign.

The situation with the \( \frac{1}{2} (110) \) *screw* dislocation is more complicated. The ions in any particular (110) row are of the same sign (see Fig. 1.12), and since displacements are parallel to the Burgers vector, motion of the screw results in displacement of charge parallel to the line. Consequently, both kinks and jogs on screw dislocations are charged, the effective charge being \( \pm e/4 \) in each case. The effective charge of the point of emergence on \{110\} and \{100\} surfaces is \( \pm e/8 \).

The reason underlying the choice of \{110\} as the principal slip plane is unclear. It has long been considered that the glide system is determined by the strengths of the electrostatic interactions within the dislocation core. This is partly supported by recent calculations. Although stable stacking faults do not exist, the core may spread on the \{110\} planes to a width of about \( 6b \), and thus consist of two fractional dislocations of Burgers vector \( \frac{1}{4} (110) \) bounding an unstable fault. Dissociation on the \{100\} planes does not occur and is small on \{111\}. It has also been suggested, however, that since the ions of the row at the bottom of the extra half-plane of the \( \frac{1}{2} (110) \) edge dislocation all have the same sign for \{100\} and \{111\} slip, but alternate in sign for \{110\} slip, interaction between edge dislocations and charged impurities may be an important factor.

### 6.5 Dislocations in Superlattices

In many substitutional solid solutions of one element, \( A \), in another, \( B \), the different species of atoms are arranged at random on the atomic positions of the lattice. At a composition, \( A_xB_{1-x} \), for example, any
given lattice point is occupied indifferently by either A or B atoms. There are some solid solutions, however, particularly near stoichiometric compositions such as \( AB, AB_2, AB_3 \), etc., in which a specific distribution of the atom species can be induced. Below a critical temperature, atoms of one kind segregate more or less completely on one set of lattice positions, leaving atoms of the other kind to the remaining positions. The resulting arrangement can be described as a lattice of A atoms interpenetrating a lattice of B atoms. A random solid solution is changed to an \textit{ordered alloy} with a superlattice. Many ordered structures exist. Two possible cubic superlattices produced by alloys of composition \( AB \) (e.g., CuZn, NiAl) and \( AB_3 \) (e.g., Cu₃Au, Ni₃Al) are shown in Figs 6.12(a) and (b). These structures are given the crystallographic identification \( B_2 \) and \( L1_2 \), respectively. In the disordered state, CuZn (\( \beta \)-brass) is body-centred cubic and Cu₃Au is face-centred cubic. In the ordered form, both are based on a simple cubic Bravais lattice, one with two atoms per unit cell, the other with four. During the nucleation and growth of ordered domains in a disordered crystal, the lattice parameter change is usually sufficiently small for the atomic planes to remain continuous. Thus, when domains meet, the A and B sublattices are either in phase, i.e. in ‘step’, with each other or out of phase. The latter condition results in an \textit{antiphase boundary (APB)}. It has a characteristic energy because the nearest-neighbour coordination of the superlattice is destroyed: typical energies are similar to those of stacking faults, i.e. \( \sim 10^{-100} \text{mJ m}^{-2} \).

Antiphase boundaries also arise in the core of dislocations in ordered alloys and this is of considerable technological significance because of their influence on the high temperature mechanical properties of these materials. At low temperature, for example, \( L1_2 \) alloys behave like the face-centred cubic metals in that the critical resolved shear stress (CRSS) is almost independent of temperature, but as the temperature is increased the CRSS actually increases, an effect known as the \textit{yield stress anomaly}. If the order–disorder transition temperature for the alloy is high enough, the yield stress reaches a peak. This occurs between 800 and 1000 K in the case of Ni₃Al and allows Ni-based superalloys that contain Ni₃Al precipitates (\( \gamma' \) phase) in a disordered matrix (\( \gamma \) phase) to be used in high temperature applications, e.g., turbine blades. The key to this property is the core structure of the dislocations responsible for slip.

In disordered \( L1_2 \) alloys, dislocation behaviour is similar to that described for face-centered cubic metals in Chapter 5, e.g., the unit slip vector is \( \frac{1}{2}(110) \). In the ordered state, \( \frac{1}{2}(110) \) vectors are not lattice translation vectors and so gliding dislocations leave behind a surface of disorder (APB). This is illustrated schematically for the \( L1_2 \) superlattice in Fig. 6.13(a), which shows the atomic arrangement in two adjacent \{111\} planes. Displacement of one layer by \( \frac{1}{2}(110) \) with respect to the other shifts \( X \) to \( Y \) and creates an APB in which A atoms occupy nearest-neighbour sites to each other. Order may be restored, however, by a second displacement \( \frac{1}{2}(110) \) which takes the atom originally \( X \) to \( Z \).
Dislocations in other crystal structures

Unlike the face-centred cubic metals, the $LI_2$ alloys are also commonly observed to slip on the $\{100\}$ planes. In this case, see Fig. 6.13(b), the shift $X$ to $Y$ leaves nearest-neighbour bonds across the APB unchanged and the second-neighbour changes are the major contribution to the APB energy, implying that the APB energy is lower on $\{100\}$ than $\{111\}$. Again, the second shift $Y$ to $Z$ restores order.

Thus, the perfect dislocation moving on either the $\{111\}$ or $\{100\}$ planes of the $LI_2$ superlattice consists of two $\frac{1}{2} \{110\}$ superpartial dislocations joined by an APB. This superdislocation has a Burgers vector $(110)$, which is a lattice translation vector. The superdislocation is similar to two Shockley partials joined by a stacking fault (section 5.3), for the spacing in equilibrium is given by a balance between the elastic repulsive force between the two superpartials and the opposing

Figure 6.13 Arrangement of atoms in two adjacent atomic planes in an $LI_2$ superlattice: (a) $(111)$ planes and (b) $(100)$ planes.
force due to the \( APB \) energy. Furthermore, if the two dislocations are each dissociated into \( \frac{1}{3}(112) \) Shockley partials in the disordered phase, as is possible in the face-centred cubic structure, they may retain this form in the ordered \( L\!I_2 \) lattice. The superdislocation would then consist of two extended dislocations connected by an \( APB \). The slip system of the ordered structure is apparently stabilised by the dislocation behaviour in the disordered state, for the Burgers vector of the superdislocation is not the shortest lattice vector, which is \( \langle 100 \rangle \) in both structures in Fig. 6.12. In the alloy of Fig. 6.12(a), which is body-centred cubic in the disordered state, the Burgers vector is \( \langle 111 \rangle \) given by two \( \frac{1}{3}(111) \) components. The requirement that dislocations in superlattices travel in pairs separated by an \( APB \) provides the strengthening mechanism referred to above. It arises because of a ‘locking’ mechanism that affects the superpartial pair in the screw orientation.

It was seen in section 5.3 that when the perfect \( \frac{1}{3}(110) \) dislocation in a face-centred cubic metal dissociates into two Shockley partials, it is restricted to glide in one \( \{111\} \) plane because the Burgers vector \( \frac{1}{3}(112) \) of a partial lies in only one \( \{111\} \) plane. The glide plane can only change when constriction occurs. The same behaviour will apply to each superpartial in an \( L\!I_2 \) alloy if there is a tendency to split into Shockley partials. The two \( \frac{1}{3}(110) \) dislocations will be separated by a \( \{111\} \) \( APB \), as shown for a screw superdislocation in Fig. 6.14(a). If one of the superpartials is not dissociated, it can glide on a \( \{100\} \) plane, as illustrated in Fig. 6.14(b). This cross slip from the \( \{111\} \) (‘octahedral’) plane to the \( \{100\} \) (‘cube’) plane may be energetically favourable if the \( APB \) energy is lower on \( \{100\} \) than \( \{111\} \), which is the implication of the simple geometrical picture of nearest-neighbour coordination described above. Furthermore, depending on the elastic anisotropy of the crystal, the force (or torque) one partial exerts on the other as a result of its stress field can enhance the stability of this arrangement.

When cross slip of a screw superdislocation onto \( \{100\} \) planes occurs, glide of the remainder of the dislocation on the \( \{111\} \) system is restricted because the Peierls stress resisting glide on \( \{100\} \) is much higher than that on \( \{111\} \). Thus, the applied stress has to be increased to maintain plastic flow. The \( \{100\} \) cross-slipped segment is known as a \textit{Kear–Wilsdorf lock}. The effect of temperature that results in the yield stress anomaly referred to above probably arises from the ease with which a \( \frac{1}{3}(110) \) superpartial can cross slip onto a \( \{100\} \) plane. If it is dissociated into two \( \frac{1}{6}(112) \) Shockley partials on a \( \{111\} \) plane, it will have to constrict, as mentioned above, and the energy barrier for that can be overcome by thermal activation. The result is increased hardening due to cross slip with increasing temperature, as seen experimentally. An example of the dislocation structure in a crystal of Ni\(_3\)(Al, Hf) deformed at 400 °C by slip on a single slip system is shown in the transmission electron microscope image taken under dark-field conditions (section 2.4) in Fig. 6.15. Not all \( L\!I_2 \) alloys exhibit anomalous hardening, however, and this probably results from a higher ratio of the \( APB \) energy on \( \{100\} \) to that on \( \{111\} \) and a lower value.
Figure 6.14  (a) Screw superpartial pair with an APB on a \{111\} plane. (b) Cross slip onto a \{100\} plane, starting the formation of a Kear–Wilsdorf lock.

Figure 6.15  Dislocation structure in Ni$_3$(Al,Hf) deformed at 400 °C. The (111) slip plane is the plane of the paper and the [101] slip direction is indicated by $\mathbf{b}$. The superdislocations are elongated in the screw direction and are bowed out in the (010) cube cross-slip plane. The edge segments are shorter. The superdislocations consist of two $\frac{1}{2}$[101] superpartials with a spacing of a few nm. (Courtesy P. Veyssiere.)

of the elastic torque. Finally, it should be noted that sessile screw configurations equivalent to the Kear–Wilsdorf lock have been found in ordered alloys with other crystal structures.
6.6 Dislocations in Covalent Crystals

The covalent bond formed by two atoms sharing electrons is strongly localized and directional, and this feature is important in determining the characteristics of dislocations. Of the many covalent crystals, the cubic structure of diamond, silicon and germanium is one of the simplest and most widely studied. Compounds such as gallium arsenide (GaAs) have the same atomic arrangement. Dislocations in these semiconductors affect both mechanical and electrical properties.

The space lattice is face-centred cubic with two atoms per lattice site, one at 0, 0, 0 and the other at $\frac{1}{4}, \frac{1}{4}, \frac{1}{4}$ (Fig. 6.16). In GaAs one atom would be gallium and the other arsenic. Each atom is tetrahedrally bonded to four nearest-neighbours, and the shortest lattice vector $\frac{1}{2} (110)$ links a second-neighbour pair. The close-packed $\{111\}$ planes have a six-fold stacking sequence $AaBbCcAaBb\ldots$ as shown in Fig. 6.17. Atoms of adjacent layers of the same letter such as $Aa$ lie directly over each other, and planar stacking faults arising from insertion or removal of such pairs do not change the tetrahedral bonding. By reference to the face-centred cubic metals, the intrinsic fault has stacking sequence $AaBbCcAaBb\ldots$ and the extrinsic fault has $AaBbCcAaBb\ldots$ Faults formed between adjacent layers of the same letter do not restore tetrahedral bonding and have high energy.

Perfect dislocations have Burgers vector $\frac{1}{2} (110)$ and slip on $\{111\}$ planes. They usually lie along $\langle110\rangle$ directions at $0^\circ$ or $60^\circ$ to the Burgers vector as a result of low core energy in those orientations (see Fig. 8.7). From consideration of dislocations formed by the cutting operations of section 3.2, two dislocation types may be distinguished. The cut may be made between layers of either different letters, e.g., $aB$, or the same letter, e.g., $bB$. Following Hirth and Lothe the dislocations produced belong to either the glide set or the shuffle set, as denoted in Fig. 6.17. Diagrammatic illustrations of the two sets are shown in Fig. 6.18. The dangling bonds formed by the free bond per atom along the core are apparent.

Dislocations of both sets are glissile and can dissociate, but the mechanism of dissociation is different for the two cases. In the glide
Dislocations in other crystal structures

Figure 6.17  (011) projection of the diamond-cubic lattice showing the stacking sequence of the (111) planes and the shuffle and glide planes defined in the text.

Figure 6.18  Perfect 60° dislocations of (a) the glide set and (b) the shuffle set. The extra half-planes are in the lower part of each drawing. (After Hirth and Lothe (1968), Theory of Dislocations, McGraw-Hill.)
Introduction to Dislocations

Figure 6.19: Two possible core configurations for the 30° partial of the glide set in silicon. (a) Dangling bonds occur at atoms CCCC. (b) Reconstructed core. (From Marklund, Phys. Stat. Sol. (b), 92, 83, 1979.)

For the glide set, the perfect dislocation dissociates into two $\frac{1}{6}(112)$ Shockley partials separated by the intrinsic stacking fault, as in the face-centred cubic metals. Dissociation of the shuffle dislocation is not so simple because of the absence of low-energy shuffle faults. It occurs by the nucleation and glide of a Shockley partial of the glide type between an adjacent pair of \{111\} layers. This results in a fault of the glide set bounded on one side by a Shockley partial and on the other by a Shockley partial and, depending on whether the glide fault is above or below the shuffle plane, a row of interstitials or vacancies. It is probable that this dislocation is less mobile than the glide-set dislocation because movement of the row of point defects within the core can only occur during slip by shuffling: hence the nomenclature. Climb, which involves point defect absorption or emission, transforms shuffle-set dislocations to glide-set dislocations, and vice versa. The stacking-fault energy in silicon and germanium is sufficiently low ($\approx 50 \text{ mJ m}^{-2}$) for dissociation to be resolved directly in the transmission electron microscope (Fig. 2.10). Computer calculations of the core energy of dislocations in model silicon crystals suggest that the energy is reduced by bond reconstruction, a process in which dangling bonds reform with others so that all atoms retain approximately tetrahedral coordination. This is shown schematically in Fig. 6.19 for the glide-set partial with $\mathbf{b} = \frac{1}{6}(112)$ at 30° to the line direction. Reconstruction occurs by the dangling bonds at CCCC rebonding in pairs along the core. The structure within the core affects the electron energy levels and, therefore, the electrical properties of crystals containing dislocations.

6.7 Dislocations in Layer Structures

There is a large group of materials with a pronounced layer-type structure which can arise in two ways. Firstly, when the binding forces between atoms in the layers are much stronger than the binding
forces between atoms in adjacent layers, as, for example, in graphite. Secondly, when the arrangement of the atoms in complex molecular structures results in the formation of two-dimensional sheets of molecules as, for example, in talc and mica. There are a number of important consequences of the layer structure. Slip occurs readily in planes parallel to the layers and is almost impossible in non-layer planes and, therefore, the dislocation arrangements and Burgers vectors are confined mainly to the layer planes. The weak binding between layers can result in stacking faults of low energy and hence unit dislocations are usually widely dissociated into partial dislocations. Figure 6.20 shows an example of dislocations in talc. In this material the unit dislocations dissociate into four component partial dislocations. The dislocations appear as ribbons lying in the layer planes, and in some circumstances the electron diffraction conditions allow all four partials to be observed.

Many studies have been made of crystals with layer structures. They are particularly convenient to study experimentally because uniformly thick specimens for transmission electron microscopy can be obtained simply by cleavage along the layer planes. The geometry of the Burgers vectors and dislocation reactions are described using the methods developed in preceding sections.
6.8 Dislocations in Polymer Crystals

The plastic deformation of crystalline and semi-crystalline polymers involves the mechanisms well established for other crystalline solids, namely dislocation glide, deformation twinning and, in some cases, martensitic transformations. As in preceding sections, dislocation behaviour can be described after making allowances for the structure. Polymer crystals are distinguished by having strong covalent bonding in the direction of the molecular chain axis, and weak van der Waals bonding in the transverse directions. This is illustrated for the orthorhombic phase of polyethylene in Fig. 6.21, which has been the subject of much investigation. As a result of the relative stiffness of the molecules, the important dislocations lie along the [001] direction. The shortest lattice vector is [001], and the screw dislocation produces chain-axis slip by gliding as a perfect dislocation on the (100), (010) and {110} planes. Transverse slip results from the glide of edge dislocations, and although the shortest lattice vectors in the transverse direction are [010] and [100], the slip vector is (110). This is because the perfect (110) dislocation can dissociate into two Shockley partials of Burgers vector approximately $\frac{1}{2}$(110) bounding a stacking fault on a {110} plane. The fault is believed to have low energy ($\leq 10 \text{ mJ m}^{-2}$), but the [001] screw is prevented from dissociating by Frank’s rule.

In bulk polyethylene, the (001) surfaces of a crystal consist of the folds formed by the molecules leaving and re-entering the crystal. The folds are aligned along certain crystallographic directions, and this influences the dislocation behaviour within the crystals by favouring slip planes containing the folds.

Further Reading

Body-centred Cubic and Hexagonal Metals


Yoo, M. H. and Wuttig (Eds.) (1994) *Twinning in Advanced Materials*, TMS.

**Other Crystal Structures**


