CHAPTER I

INTRODUCTION

Galvanomagnetic phenomena in ferromagnetic metals and alloys have aroused great interest in a large number of workers for a period which encompasses the last four decades. The past decade, in particular, has witnessed greatly expanded activity, mostly in the nature of refinements of the basic theory and application to specific examples, with particular attention to special cases of band structure and scattering mechanisms. The need to pursue such new lines of thinking has arisen from the fact that there still exists a considerable speculation regarding the possible source of the peculiarities in the transport properties of such materials. The anomalous magnetoresistivity, its unusual dependence on temperature, impurity content and applied magnetic field, the failure of Kohler's rule in certain temperature ranges, the increase in the electrical resistivity faster than the first power of temperature near the Curie point and the discontinuity in its temperature derivative at the Curie temperature, the deviations from Matthiessen's rule at low temperatures are only a few amongst many of the special features of ferromagnetics. The conventional theories of metallic conduction (Bloch, 1928, 1930) and magnetoresistivity (two-band model) fail to provide a satisfactory answer for these peculiarities as can be seen from the example that the resistivity of ferromagnetic as well as non-ferromagnetic transition metals deviates from the well-known Bloch-Gruneisen T^5 law at low temperatures, where,

instead a T^2 behaviour is approached. These observations point
to the fact that in the earlier theories no account has been
taken of the extra complexity introduced by the influence of the
exchange interactions, which order the spins in a ferromagnet.
In order to overcome these inadequacies in the previous models,
many of the subsequent models have been conceived. In the sec-
tions that follow an appraisal of the different theories has been
made on the basis of the available experimental data on various
ferromagnetic metals and alloys. Although quantum theory of trans-
port constitutes a major portion of the theoretical effort in this
field in recent years, almost invariably it is found that the semi-
classical theory (Boltzmann transport equation) forms an excellent
and reliable approximation, and has, moreover, the great virtue of
lending itself to vivid physical interpretations. Thus, the next
section is devoted to the formulation and the general solution of
Boltzmann transport equation.

(i) <u>Formulation and solution of the transport equation</u> :

It is customary to describe the assembly of conducting
electrons in a metal by means of a distribution function $f(\underline{K},\underline{r},t)$,
defined such that the number of electrons in the six-dimensional
volume element $d\underline{K} \, d\underline{r}$ at time t is given by $(1/4\pi^3) \, f(\underline{K},\underline{r},t)$
$d\underline{K} \, d\underline{r}$. In principle, the electronic transport properties of
a conductor are completely specified once $f(\underline{K},\underline{r},t)$ is known.
Boltzmann transport equation, in its fundamental form, can be
written as

$$\left(\frac{\partial f}{\partial t}\right)_d + \left(\frac{\partial f}{\partial t}\right)_c = 0 \qquad \qquad \cdots (1)$$

The first term in equation (1) represents two processes which bring about changes in the distribution function, namely, (i) changes due to the drift of electrons under the action of externally applied fields and (ii) changes in distribution due to thermal gradients, whereas the second term takes account of the changes in f due to collisions of electrons within themselves and with imperfections in the lattice. Equation (1) represents a steady state as distinguished from the equilibrium state, denoted by $f_0(\epsilon)$, which holds when the external fields and temperature gradients are absent.

If the drift, diffusion and collision (Nordheim, 1931) terms in equation (1) are written in detail, Boltzmann transport equation assumes its usual form

$$\frac{e}{\hbar}\left[\underline{E} + \frac{1}{c}\,\underline{v}\times\underline{H}\right]\cdot\nabla_{\underline{k}}\,f \;+\; \underline{v}\cdot\nabla_{\underline{r}}\,f$$

$$= \int\left\{Q(\underline{k}',\underline{k})\,f(\underline{k}')\left[1-f(\underline{k})\right] - Q(\underline{k},\underline{k}')\,f(\underline{k})\left[1-f(\underline{k}')\right]\right\}d\underline{k}' \qquad \dots (2)$$

where $Q(\underline{k}', \underline{k})$ is the probability per unit time that an electron makes a transition from the state $\underline{k}'$ to the state $\underline{k}$ and the weight factors $f(\underline{k}')$ and $\left[1 - f(\underline{k})\right]$ ensure that the state $\underline{k}'$ is occupied by an electron while the state $\underline{k}$ is unoccupied ; thereby taking proper account of Pauli exclusion principle. Similar interpretation applies to the second term in the integrand of right-hand side.

The solution of the integro-differential equation (2) can be obtained in closed form only when the collision integral is expressed in terms of a relaxation time τ i.e., when $(\partial f/\partial t)_c =$

$-(f-f_0)/\tau$. The physical significance of τ is that of a time constant in the exponential approach of a disturbed distribution f to the equilibrium distribution f_0 when the external interactions are suddenly switched off.

The relaxation time approximation is not always a useful concept particularly when the mechanism of scattering is complex. A meaningful relaxation time $\tau(\underline{K})$ can be defined whenever the energy change of an electron per collision is small compared to kT, the thermal energy. In other words, a relaxation time can be defined unambiguously when applied to scattering of electrons by acoustical phonons or impurities in rather impure metals and pure metals at temperatures higher than their Debye temperature. In the temperature range in which this approximation fails, transition probabilities are to be calculated explicitly by a detailed consideration of the scattering processes. The general function consists in expanding $f(\underline{K}, \underline{r})$ about the equilibrium distribution function $f_0(\underline{K}, \underline{r})$ as

$$f(\underline{K}, \underline{Y}) = f_0(\epsilon) + f_1(\underline{K}, \underline{Y})$$

$$= f_0(\epsilon) - \phi(\underline{K}, \underline{Y}) \frac{\partial f_0}{\partial \epsilon} \qquad \ldots (3)$$

Such an expression indicates that the electric field simply shifts the distribution in $\underline{K}$ space unaltered in shape and the current density has a non-zero value only over a small region near the Fermi surface where $\partial f_0/\partial \epsilon$ is finite. ϕ is thus a measure of the deviation caused by external fields in the distribution function from its equilibrium value. Hence, the problem reduced to finding out the solution for ϕ. In the absence of thermal

gradients, equation (2) can be rewritten in terms of ϕ (in relaxation time approximation) as

$$\phi = e \tau \underline{E} \cdot \underline{V} + \frac{e\tau}{c\hbar^2}\left(\underline{H} \cdot \underline{\Omega} \phi\right) \qquad \dots (4)$$

where $\quad \underline{\Omega} = \nabla_{\underline{K}} \epsilon \times \nabla_{\underline{K}}$

Although a formal solution of (4) can be written in closed form, it is so cumbersome that it has rarely been employed for computations. Instead, physically meaningful solutions can be obtained if the following two situations are considered :

(1) $H \neq 0$, $\epsilon = \hbar^2 K^2/2m^*$, i.e. the Fermi surface is spherical and (2) $H \neq 0$, ϵ an arbitrary, well-behaved function of $\underline{K}$.

In the first case, the simplification arises from the fact that now $\underline{V} = \hbar \underline{K}/m^*$ and equation (4) yields the solution

$$\phi = e\tau\underline{V}\cdot\left[\underline{E} - \alpha\tau\underline{H}\times\underline{E} + (\alpha\tau)^2\underline{H}(\underline{H}\cdot\underline{E})\right]/\left[1 + (\alpha\tau H)^2\right] \qquad \dots (5)$$

where $\quad \alpha = \dfrac{e}{m^*c}$.

It is evident from equation (5) that if $\underline{H}$ and $\underline{E}$ are colinear, it reduces to $\phi = e\tau\underline{E}\cdot\underline{V}$ which is the solution of equation (4) in the case $H = 0$. Thus, it follows that under these conditions the magnetic field has no influence whatever on the transport properties of the electron gas. Alternatively, a quasi-free-electron gas cannot display longitudinal magnetoresistance.

In the second case, the solution of equation (4) lends itself to computation only in the low-field limit, $\alpha\tau H \ll 1$.