This book has been written using gaussian cgs units throughout, except for a few formulas that are explicitly quoted in practical units for convenience in applications. This choice conforms to that in the first edition and to that used in much of the physics literature on which the book is based. Moreover, because of the central role of \( \mathbf{B} \) and \( \mathbf{H} \) in superconductivity, it is especially convenient to have \( \mathbf{B} = \mathbf{H} \) in vacuum and the natural form \( \mathbf{E} = (\mathbf{v}/c) \times \mathbf{B} \) for the electric field associated with a moving flux density.

For readers who are more comfortable with SI or mksa units, we reproduce in Table A.1 a version of the convenient tabular conversion guide given in the appendix of J. D. Jackson, Classical Electrodynamics, Wiley, New York, 1975, p. 819. To convert any formula from gaussian to SI units, follow these rules: Symbols for mass, length, time, force, and other quantities that are not specifically of an electromagnetic nature are unchanged. Symbols for electromagnetic quantities listed under “gaussian” in Table A.1 are replaced on both sides of the equation by the corresponding symbols listed under “SI.” The reverse transformation can also be made. To illustrate this procedure, we consider a few important examples.

The flux quantum \( \Phi_0 = hc/2e \) in gaussian units. Following Table A.1, the left-hand side of this equation becomes \( \sqrt{4\pi/\mu_0} \Phi_0 \) and the right-hand side becomes \( h(1/\sqrt{\mu_0 \varepsilon_0})(\sqrt{4\pi \varepsilon_0/2e}) \). After canceling common factors, one obtains \( \Phi_0 = h/2e \), which defines the flux quantum in SI units. Similar manipulations applied to the GL relation (4.20) relating the flux quantum to the product of \( H_c \xi \lambda \) leave the form of the equation unchanged, except that the notation \( H_c \) must be replaced by \( B_c \) because \( \Phi_0 \) is a quantum of magnetic flux or induction, not
TABLE A.1
Conversion table for electromagnetic formulas

<table>
<thead>
<tr>
<th>Quantity</th>
<th>Gaussian</th>
<th>SI (mkSa)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Velocity of light</td>
<td>$c$</td>
<td>$\frac{1}{\sqrt{\mu_0 \varepsilon_0}}$</td>
</tr>
<tr>
<td>Magnetic induction or flux density</td>
<td>$B$</td>
<td>$\sqrt{\frac{4\pi}{\mu_0}} B$</td>
</tr>
<tr>
<td>Magnetic field</td>
<td>$H$</td>
<td>$\sqrt{4\pi \mu_0} H$</td>
</tr>
<tr>
<td>Magnetization</td>
<td>$M$</td>
<td>$\sqrt{\frac{\mu_0}{4\pi}} M$</td>
</tr>
<tr>
<td>Charge density (or charge, current, current density, polarization)</td>
<td>$\rho$ (or $Q, I, J, P$)</td>
<td>$\frac{1}{\sqrt{4\pi \varepsilon_0}} \rho$ (or $Q, I, J, P$)</td>
</tr>
<tr>
<td>Electric field (or potential, voltage)</td>
<td>$E$ (or $\phi, V$)</td>
<td>$\sqrt{4\pi \varepsilon_0} E$ (or $\phi, V$)</td>
</tr>
<tr>
<td>Displacement</td>
<td>$D$</td>
<td>$\sqrt{\frac{4\pi}{\varepsilon_0}} D$</td>
</tr>
<tr>
<td>Conductivity</td>
<td>$\sigma$</td>
<td>$\frac{\sigma}{4\pi \varepsilon_0}$</td>
</tr>
<tr>
<td>Resistance (or impedance)</td>
<td>$R$ (or $Z$)</td>
<td>$4\pi \varepsilon_0 R$ (or $Z$)</td>
</tr>
<tr>
<td>Inductance</td>
<td>$L$</td>
<td>$4\pi \varepsilon_0 L$</td>
</tr>
<tr>
<td>Capacitance</td>
<td>$C$</td>
<td>$\frac{C}{4\pi \varepsilon_0}$</td>
</tr>
<tr>
<td>Permeability</td>
<td>$\mu$</td>
<td>$\frac{\mu}{\mu_0}$</td>
</tr>
<tr>
<td>Dielectric constant</td>
<td>$\varepsilon$</td>
<td>$\frac{\varepsilon}{\varepsilon_0}$</td>
</tr>
</tbody>
</table>


field. Similarly, the condensation energy per unit volume $F_s = F_0 = H^2_c / 8\pi$ becomes $\mu_0 H^2_c / 2$ or $B^2_c / 2 \mu_0$ in SI units, and the London penetration depth $\lambda = \sqrt{m c^2 / 4\pi n_0 e^2}$ becomes $\sqrt{m / \mu_0 n_0 e^2}$. The definition of the coherence length $\xi$ is unchanged because it does not involve electromagnetic quantities, but only $\hbar$, $m$, and the GL coefficient $\alpha$, all of which are unchanged.
NOTATION FOR MAGNETIC FIELDS. Because of the ubiquity of magnetic fields in the subject of superconductivity, special notational conventions are common to simplify the discussion. We follow the convention of de Gennes (and others) and use \( \mathbf{h}(\mathbf{r}) \) to denote the local value of the magnetic induction or flux density, which typically varies on the scale of the penetration depth \( \lambda \). We reserve the use of \( \mathbf{B} \) to denote the value of \( \mathbf{h} \) averaged over such microscopic lengths but still capable of varying smoothly over the macroscopic dimensions of the sample.

In normal metal or vacuum, of course, there is no microscopic variation of \( \mathbf{h} \) (we neglect the Landau diamagnetism and Pauli paramagnetism), so \( \mathbf{B} = \mathbf{h} \). In these cases, \( \mathbf{B} = \mathbf{H} \), so all three symbols denote equal quantities, and may be used interchangably.

In the Meissner state of a massive superconductor, \( \mathbf{h} \) is reduced to zero within a penetration depth \( \lambda \) of the surface by supercurrents in the skin layer, as described by the Maxwell equation

\[
\text{curl} \mathbf{h} = \frac{4\pi J_{\text{total}}}{c} \tag{A2.1}
\]

Hence, \( \mathbf{B} = \mathbf{h} = 0 \) deep inside. On the other hand, \( \mathbf{H} \) is governed by the Maxwell equation

\[
\text{curl} \mathbf{H} = \frac{4\pi J_{\text{ext}}}{c} \tag{A2.2}
\]
where $J_{ext}$ represents a nonequilibrium current and excludes currents arising from the equilibrium response of the medium, such as those in the penetration depth described by the London equations. Hence, curl $\mathbf{H} = 0$, and the tangential component $H_t$ is constant through the skin depth, retaining the value of $H_t(=B_t = h_t)$ found outside the sample. If the sample is ellipsoidal, $H$ inside is uniform and everywhere equals the equatorial value of $H_t$. This $H_t$ will in general exceed the uniform applied field $H_0$ by a factor $(1 - \eta)^{-1}$, where $\eta$ is the shape-dependent demagnetizing factor of the sample. [See the discussion associated with (2.20).]

In the intermediate state of a type I superconductor, which is reached when $H_t = H_c$, the magnitude of $h$ varies continuously (on the scale of $\lambda$) between $H_c$ in the normal lamina and zero in the superconducting ones. $B$ is the average of this $h$ over the laminar structure, and it is constant within an ellipsoidal sample. The magnitude of $\mathbf{H}$ must be $H_c$ for coexistence of superconducting and normal regions to be possible. These interrelations are illustrated in Fig. 2.3.

In the mixed state of a type II superconductor in a magnetic field above $H_{cm}$, $h$ varies on the microscopic scale of the vortex structure, whereas $B$ is the average of $h$ over the structure. In the ideal equilibrium case, $H$ is again everywhere equal to $H_t$ at the equatorial surface. In the presence of transport currents and of disequilibrium due to flux pinning, $H$ will vary because curl $\mathbf{H}$ is no longer zero. This situation is discussed more fully in connection with (5.48) in the text.

**NOTATION FOR ELECTRIC FIELDS.** For notational symmetry, we also can define a microscopically varying electric field $\mathbf{e}(\mathbf{r})$, whose macroscopic average is $\mathbf{E}$. But because curl $\mathbf{e} = -(1/c)\partial \mathbf{h}/\partial t$, $\mathbf{e}$ is uniform in static situations, and it must be zero in equilibrium. Thus, the distinction between $\mathbf{e}$ and $\mathbf{E}$ arises less frequently than the distinction between $\mathbf{h}$ and $\mathbf{B}$ (which can result from equilibrium supercurrents); consequently, we have normally used $\mathbf{E}$ for both to avoid confusion between $\mathbf{e}$ and the electronic charge $e$. The notation $\mathbf{e}$ is introduced only to describe the electric field distribution about a moving vortex, where the macroscopic average $\mathbf{E}$ is quite different from $\mathbf{e}$ and gives the physically important resistive voltage.

**SIGN OF THE ELECTRON CHARGE.** Fortunately, the sign convention for the electronic charge is often immaterial since only its square enters into such quantities as the penetration depth or normal conductivity. Also, in the definition of quantities such as the flux quantum $\Phi_0 = hc/2e$ or the Josephson frequency $2eV/h$, it is convenient as well as conventional to take $e$ to be a magnitude. However, when the sign of the charge matters, we have chosen to follow de Gennes in adopting the convention that $e$ is the charge of the electron, including its sign; i.e., $e = -|e|$. This means that a metal island with $n$ excess electrons on it has a charge $ne$, not $-ne$, and the supercurrent density is simply $n_e \mathbf{v}_s$, rather than $n_e(-e)\mathbf{v}_s$. This convention simplifies many expressions, and seems more physical. It has the disadvantage of being contrary to the convention used in such popular textbooks as *Introduction to Solid State Physics* by Kittel and *Solid State Physics* by Ashcroft and Mermin, in which the charge of the electron is written as $(-e)$, so that $e$ always refers to a magnitude.
A convenient technique for obtaining an exact expression for the penetrating magnetic field, and hence for the penetration depth, is to apply Fourier analysis to $\mathbf{J}$ and $\mathbf{A}$, and to use (3.101) to obtain a self-consistent solution. Only a one-dimensional Fourier analysis is required since $J_x$ and $A_x$ are functions only of $z$ for the penetration of a magnetic field $B_z$ parallel to a planar surface. Some care is needed in handling the surface, however, since our expressions for the response function $K(q)$ are valid only in an infinite medium. This problem is handled by the mathematical artifice of introducing externally supplied source currents in the interior of the infinite medium to simulate the field applied at a surface.

Consider, e.g., the case in which electrons are assumed to be specularly reflected at the surface. If one introduces a current sheet

$$J_{x,\text{ext}} = -\frac{c}{2\pi} B_0 \delta(z)$$

(A3.1)

this introduces a discontinuity $2B_0$ in $h_y$. This can be taken symmetric about zero, so that $h_y$ switches from $-B_0$ to $+B_0$. Now when the superconductive medium is introduced, its diamagnetic currents will screen out these fields in a length $\lambda$ (to be determined). Note that electrons passing through this plane at $z = 0$ without
scattering have had a past exposure along their trajectory to a vector potential exactly the same as that seen by electrons specularly reflected at the surface in the actual case since $\mathbf{A}(-z) = \mathbf{A}(z)$. (See Fig. A3.1.) Thus, the net supercurrent induced in them should also be the same, and the simulation should be effective.

Having replaced the surface by a current sheet in an infinite medium, we now may proceed to use the response function $K(q)$ worked out for that case. We first note that

$$\nabla^2 \mathbf{A} = -\text{curl} \ \text{curl} \ \mathbf{A} = -\text{curl} \ \mathbf{h} = -\frac{4\pi}{c} \mathbf{J}_{\text{total}} = -\frac{4\pi}{c} (\mathbf{J}_{\text{ext}} + \mathbf{J}_{\text{med}})$$

For the $q$th Fourier component, this becomes

$$q^2 a(q) = \frac{4\pi}{c} \mathbf{J}_{\text{ext}}(q) - K(q) a(q)$$

Solving for $a(q)$, we have the general result

$$a(q) = \frac{(4\pi/c)\mathbf{J}_{\text{ext}}(q)}{K(q) + q^2} \quad (A3.2)$$

For the current sheet (A3.1), $\mathbf{J}_{\text{ext}}(q) = -cB_0/4\pi^2$, and we drop the vector notation since $\mathbf{J}$ and $\mathbf{A}$ have only an $x$ component. Thus,

$$a(q) = \frac{-B_0/\pi}{K(q) + q^2}$$

**FIGURE A3.1**

Simulation of surface with specular reflection by source-current sheet. (a) Magnetic field in normal (dashed) and superconducting (solid) states. (b) Vector potential in normal (dashed) and superconducting (solid) states. London gauge is used in superconducting state. (c) Electron trajectories. The solid curve shows trajectory with specular reflection; the dashed parts show extensions into the other half-space, with current-sheet simulation.
We are more interested in \( h = \text{curl} \mathbf{A} \), so that \( h_y(q) = i q a(q) \). Integrating over all the Fourier components, we obtain

\[
    h(z) = \frac{B_0}{\pi r} \int_{-\infty}^{\infty} \frac{q e^{i q z}}{K(q) + q^2} \, dq = \frac{2B_0}{\pi} \int_{0}^{\infty} \frac{q \sin qz \, dq}{K(q) + q^2} \tag{A3.3}
\]

For any \( K(q) \), (A3.3) gives the true dependence of \( h \) on \( z \), which will not be exactly exponential unless \( K(q) = \) constant, as in the London theory. For example, the \( h(z) \) computed with the \( K(q) \) for either the Pippard or BCS theory actually changes sign deep in the interior, where \( |h(z)| \ll B_0 \).

To get the penetration depth, as usually defined, we integrate (A3.3):

\[
    \lambda = B_0^{-1} \int_{0}^{\infty} h(z) \, dz = \frac{2}{\pi} \int_{0}^{\infty} \frac{q \sin qz \, dq \, dz}{K(q) + q^2}
\]

or

\[
    \lambda_{\text{spec}} = \frac{2}{\pi} \int_{0}^{\infty} \frac{dq}{K(q) + q^2} \tag{A3.4}
\]

(In carrying out the integration on \( z \), one can replace \( \int_{0}^{\infty} q \sin qz \, dq \) by its average value, unity, since as \( Z \to \infty \), the oscillatory part effectively averages to zero in the subsequent integration over \( q \).)

Given (A3.4), we can compute \( \lambda_{\text{spec}} \) for any model of superconductivity which determines a \( K(q) \). For example, in the London theory, \( K(q) = 1/\lambda_L^2 \). Then

\[
    \lambda_{\text{London, spec}} = \frac{2}{\pi} \int_{0}^{\infty} \frac{dq}{\lambda_L^{-2} + q^2} = \lambda_L \tag{A3.5}
\]

In the Pippard theory, one has

\[
    K_p(q) = \frac{1}{\lambda_L^2 \xi_0} \left\{ \frac{3}{2(q \xi)^3} \left[ (1 + q^2 \xi^2) \tan^{-1} q \xi - q \xi \right] \right\} \tag{A3.6}
\]

This is found from (3.117) with \( J_p(R, T) = e^{-R/\xi} \) by using the general relation (3.106). If instead one approximates the BCS kernel even more closely by \( J(R, T) \approx J(0, T) \exp \left[ -J(0, T) R/\xi_0 \right] \), as discussed in the argument leading to (3.123), the effect is simply to replace \( \xi_0 \) by \( \xi_0' = \xi_0/J(0, T) \) everywhere in (A3.6), including in the definition (3.121) of \( \xi \). As remarked in Chap. 3, these rather convenient, generalized Pippard forms provide quite a serviceable approximation to the exact numerical results of BCS. However, even with the analytic expression (A3.6) for \( K(q) \), numerical integration is required to compute the penetration depth by using (A3.4).

In order to avoid numerical calculations, considerable attention has been given to two limiting cases in which analytic results can be obtained, even though the true situation usually lies in between.
The local approximation replaces $K(q)$ for all $q$ by $K(0)$, a constant, thus reducing the problem to the London form, but in general with a modified penetration depth. Using the generalized Pippard approximation

$$K(0, T) = \lambda_L^{-2} \left[ 1 + \frac{\xi_0}{J(0, T) \ell} \right]^{-1} \quad \text{(A3.7)}$$

one finds

$$\lambda(T) = \lambda_L(T) \left[ 1 + \frac{\xi_0}{J(0, T) \ell} \right]^{-1/2} \quad \text{(A3.8)}$$

as anticipated in (3.123). This approximation is reasonably well justified in dirty superconductors [if $\ell < \lambda(T)$], in high-temperature superconductors, and even in pure classic superconductors very near $T_c$ where $\xi_0 < \lambda(T)$.

The other approximation is the extreme anomalous limit, in which $K(q)$ is replaced for all $q$ values by its asymptotic form for $q \to \infty$, where $K(q) \sim 1/q$. This approximation is reasonably well justified if $\lambda_L \ll \xi_0$ because then the dominant contribution to (A3.4) will come from the $q$ values in which this asymptotic form is valid. Figure A3.2 illustrates the two different approximations to $K(q)$. Since both approximations exceed the true $K(q)$ for some $q$ and never err in the other direction, both will lead to lower bounds to the true value for $\lambda$.

Let us now carry out the calculation in the extreme anomalous limit. For complete generality, we write

$$K(q) = \frac{a}{q}$$

where in the Pippard theory $a = 3\pi/4\lambda_L^2 \xi_0$, whereas in either the BCS theory or the generalized Pippard theory, $\xi_0$ is replaced by $\xi_0'$ so that $a$ is increased by a factor of $J(0, T)$. If we introduce the standard notation $\lambda_\infty$ for the value of $\lambda$ in this limit, (A3.4) becomes

$$\lambda_{\infty, \text{spec}} = \frac{2}{\pi} \int_0^\infty dq \frac{dq}{(a/q) + q^2}$$

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{figA3.2}
\caption{Schematic comparison of local and extreme anomalous approximations to the exact nonlocal response function $K(q)$.}
\end{figure}
Making a change of variable to $x = q^2/a$, we obtain

$$\lambda_{\infty, \text{spec}} = \frac{2}{3\pi a^{1/3}} \int_{0}^{\infty} \frac{x^{-1/3} \, dx}{1 + x} = \frac{4}{3\sqrt{3} a^{1/3}}$$

Inserting the value for $a$, we have

$$\lambda_{\infty, \text{spec}} = \frac{8}{9} \frac{3^{1/6}}{(2\pi)^{1/3}} (\lambda_{L}^{2} \xi_{0}^{2})^{1/3} = 0.58 (\lambda_{L}^{2} \xi_{0}^{2})^{1/3} \quad (A3.9)$$

which has exactly the form anticipated in (3.130) by an elementary argument. Since the BCS correction factor $(\xi_{0}^{2}/\xi_{0})^{1/3} = [J(0, T)]^{-1/3}$ to the simple Pippard form varies smoothly from 1 at $T = 0$ to 0.91 at $T_{c}$, it has little effect on the behavior of the result.

If the surface scattering is taken as diffuse instead of specular, formulas are obtained that differ only in detail from those given above. In this case, the prescription for handling the surface is simply to cut off the integration over $r'$ at the surface in the coordinate-space form (3.117) of the response function. The physical reasoning is that electrons coming to $r$ from the surface do so with no memory of any previous exposure to the field. When this prescription is transcribed into Fourier transform language, it turns out that (A3.4) is replaced by

$$\lambda_{\text{diff}} = \frac{\pi}{\int_{0}^{\infty} \ln[1 + K(q)/q^{2}] \, dq} \quad (A3.10)$$

Although this looks quite different from (A3.4), it actually gives exactly the same result in the local approximation, and for $\lambda_{\infty}$ it differs from (A3.9) only in that the factor of $8/9$ is missing. Thus, there is little difference in the results for these two different limiting assumptions about the surface scattering.
BIBLIOGRAPHY

While in no way exhaustive, the following list includes many of the standard references for further reading, with a brief indication of their individual features. Within each broad category, these are listed in reverse chronological order. However, it should be noted that many of the older references are classics which are still widely used.

Monographs


Schrieffer, J. R.: *Theory of Superconductivity*, W. A. Benjamin, New York (1964). A good account of the theory by one of its founders, including Green's function topics, which are not included in this book.


Collections


Parks, R. D. (ed.): *Superconductivity*, two vols., Dekker, New York (1969); reissued by the publisher in 1992. This two-volume treatise, with chapters written by two dozen distinguished authors on their special areas of interest, is the most comprehensive available treatment of the subject as it stood in 1968.

**Review articles**


A review of the status of detailed calculations.


**Applied Superconductivity Conference proceedings**

The proceedings of the biennial Applied Superconductivity Conferences provide an ongoing source of up-to-date surveys of the state of the field. Some of the recent ones are found in:


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