

The constants are different for all values of T , S and $I \leq 2$. Only a_1 , b_1 , and c_1 are fixed in such a way that we obtain the OPEP-potential for large distances. For $I > 2$, Reid uses the OPEP-potential.

4.3 Microscopic Effective Interactions

The bare nucleon-nucleon force has—as we have already seen in the preceding section—certain features that are rather difficult to handle in practice.

There is, for instance, the hard core (or at least the very repulsive core), which would make some of the usual concepts of nuclear many-body physics extremely complicated if not inapplicable (as in the Hartree-Fock case; see Chap. 5).^{*} This comes from the infiniteness of the matrix elements of a force with a hard core. In these theories, a way out of this situation is to use, in place of the bare interaction, a so-called effective interaction, which is itself an infinite sum of scattering processes of two nucleons in the nuclear medium. The bare interaction is then simply the Born term of this series. The object of this procedure is twofold: First, in re-summing the series one gets rid of the hard core problem, since the new interaction is well-behaved at short distances. Second, we can show that in replacing the bare interaction by its effective counterpart we have at the same time consistently summed up more of the many-body effects than if one had taken just the bare interaction.

The main fields of application are: (i) the ground state properties of nuclei, where the scattering of two nucleons within the nuclear medium has to be considered; (ii) the forces between the so-called valence nucleons; and (iii) effective forces between “particles” and “holes.” There are also effective three-body forces which we will ignore in this section.

4.3.1 Brückner's G -Matrix and Bethe Goldstone Equation[†]

One of the most important effective interactions in nuclear physics is the so-called Brückner G -matrix [Br 55, Da 67, and references therein]. It is, for two nucleons in the nuclear medium—in a sense yet to be specified—the analogue of the scattering matrix for two nucleons in free space.

We therefore start our considerations with the Lippmann-Schwinger equation for the scattering matrix (T -matrix; see Fig. 4.1) of two particles

*In this section, we must quite often anticipate theories and methods which are only treated later in this book. This is contrary to our usual strategy, which is to avoid this situation as much as possible. As the reader will notice, however, the microscopic theory of effective interactions is not in a very satisfactory state, so we prefer to give a short survey here together with the description of phenomenological forces, rather than devote an extra chapter to it later. (See also Appendix F.)

[†]The discussion in this section is partially based on Gomes, Walecka, and Weisskopf [GWW 58] and the textbook of Fetter and Walecka [FW 71].

(Messiah [Me 61] Chap. XIX, Sec. 14):

$$T_{\mathbf{k}_1 \mathbf{k}_2, \mathbf{k}'_1 \mathbf{k}'_2}^E = \bar{v}_{\mathbf{k}_1 \mathbf{k}_2, \mathbf{k}'_1 \mathbf{k}'_2} + \frac{1}{2} \sum_{\mathbf{p}_1 \mathbf{p}_2} \bar{v}_{\mathbf{k}_1 \mathbf{k}_2, \mathbf{p}_1 \mathbf{p}_2} \frac{1}{E - (\mathbf{p}_1^2/2m) - (\mathbf{p}_2^2/2m) + i\eta} T_{\mathbf{p}_1 \mathbf{p}_2, \mathbf{k}'_1 \mathbf{k}'_2}^E, \quad (4.38)$$

where $\mathbf{k}_1, \mathbf{k}_2$ and $\mathbf{k}'_1, \mathbf{k}'_2$ are the momenta of the incoming and outgoing particles, respectively, and E is the total scattering energy.

If we consider the scattering of two nucleons within a nuclear medium we can show (this is derived in Sec. F.4) that it makes sense to define a scattering matrix G^E analogous to that for free particles. The changes to be made for nucleons in a nucleus are almost obvious: the plane wave indices have to be changed to shell model indices, the kinetic single-particle energies figuring in the denominator of the r.h.s. of Eq. (4.38) have to be replaced by the corresponding shell model energies, and the sum over the intermediate states has to be restricted so that it does not include states below the Fermi surface. This latter feature comes from the fact that two nucleons below the Fermi surface can only scatter into states above the Fermi surface, because all other levels are occupied and are thus excluded by the Pauli principle. Therefore, we get the following equation for the G -matrix, which is usually known under the name Bethe–Goldstone equation [BG 57] (for its mathematical derivation, see Sec. F.4).

$$G_{ab, cd}^E = \bar{v}_{ab, cd} + \frac{1}{2} \sum_{\substack{m, n \\ > \epsilon_F}} \bar{v}_{ab, mn} \frac{1}{E - \epsilon_m - \epsilon_n + i\eta} G_{mn, cd}^E, \quad (4.39)$$

where ab, \dots, mn are shell-model indices and ϵ_F is the Fermi energy. This equation is usually represented graphically in an obvious way, as shown in Fig. 4.2. Two lines connecting two interactions represent the “propagator” $1/(E - \epsilon_n - \epsilon_m)$. (More will be explained about graphs in Chap. 8 and Appendix F.) For $E < \epsilon_F$, we can ignore the $i\eta$ in the denominator of (4.39), and in this case the G -matrix is obviously Hermitian as can be checked immediately by iterating Eq. (4.39). Equation (4.39) is also often

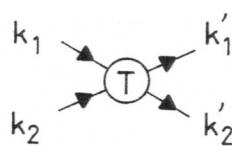


Figure 4.1. Graphical representation of the T -matrix.

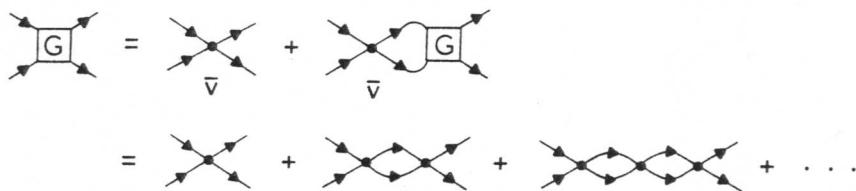


Figure 4.2. Graphical representation of the Bethe–Goldstone equation.

written in the operator form:

$$G = \bar{v} + \bar{v} \frac{Q_F}{E - H_0} G, \quad (4.40)$$

where H_0 is the shell model Hamiltonian and

$$Q_F = \sum_{\substack{m < n \\ > \epsilon_F}} |mn\rangle \langle mn| \quad (4.41)$$

is a projection operator excluding all occupied states.

In a very qualitative way, we can see from Eq. (4.40) how it can happen that G stays finite for points where \bar{v} is infinite. Solving Eq. (4.40) formally yields

$$G = \frac{\bar{v}}{1 - \bar{v} Q_F / (E - H_0)}. \quad (4.42)$$

If \bar{v} tends to infinity, G may stay finite. This is, of course, only a very crude argument and we shall in a definite example show in detail how this occurs. Before this, we have to discuss, however, some general features of Eq. (4.39). Despite the fact that Eq. (4.39) is formally very similar to Eq. (4.38), there are certain essential differences concerning, for instance, the boundary condition in the case of $E < \epsilon_F$ of the wave function defined in analogy to the scattered wave of two free particles by [Me 61, Chap. XIX 10]:

$$\begin{aligned} |\psi_{ab}\rangle &\equiv \bar{v}^{-1} G |ab\rangle = \left(1 + \frac{Q_F}{E - H_0} G \right) |ab\rangle \\ &= |ab\rangle + \frac{Q_F}{E - H_0} \bar{v} |\psi_{ab}\rangle, \end{aligned} \quad (4.43)$$

with

$$\langle \mathbf{r}_1 \mathbf{r}_2 | ab \rangle = \frac{1}{\sqrt{2}} (\varphi_a(\mathbf{r}_1) \varphi_b(\mathbf{r}_2) - \varphi_a(\mathbf{r}_2) \varphi_b(\mathbf{r}_1)) \quad (4.44)$$

and

$$H_0 |ab\rangle = (\epsilon_a + \epsilon_b) |ab\rangle. \quad (4.45)$$

The wave equation (4.43) formally resembles an equation for a scattered wave. However, contrary to the real scattering case, where the second term on the r.h.s. of (4.43) gives the outgoing scattered part of the wave function and thus does not vanish as the relative distance of the two particles $|\mathbf{r}_1 - \mathbf{r}_2|$ goes to infinity, in the present case of $E < \epsilon_F$, this term vanishes as $|\mathbf{r}_1 - \mathbf{r}_2|$ goes to infinity. This comes from the fact that for a real scattering process the T matrix of Eq. (4.38) enters Eq. (4.43) "on the energy shell," that is, the absolute value of the relative momenta of the two particles before and after the scattering process have to be the same; also, E must correspond to this value of the relative momentum. As a consequence, the corresponding energy denominator of Eq. (4.43) for a real scattering

process can be zero and introduce a singularity. In the present case this can never happen, since we suppose $E < \epsilon_F$ and therefore the second expression on the r.h.s. of Eq. (4.39) is never "on the energy shell," which is why it vanishes for large values of $|\mathbf{r}_1 - \mathbf{r}_2|$. A derivation is presented in a review article by Day [Da 67] which, however, would be too much of a digression to be repeated here. We have thus:

$$|\psi_{ab}^E\rangle \xrightarrow{|\mathbf{r}_1 - \mathbf{r}_2| \rightarrow \infty} |ab\rangle. \quad (4.46)$$

From Eqs. (4.43) and (4.41) we also immediately get the normalization condition for $a, b < \epsilon_F$:

$$\langle ij | \psi_{ij} \rangle = 1, \quad i, j < \epsilon_F. \quad (4.47)$$

It is also instructive to expand the correlated pair function $|\psi_{ij}\rangle$ in an uncorrelated basis. From (4.43) we have:

$$|\psi_{ij}\rangle = |ij\rangle + \frac{1}{2} \sum_{mn} C_{mn}^{ij} |mn\rangle, \quad (4.48)$$

that is, the correlated function contains, in addition to $|ij\rangle$, components above the Fermi level. It turns out that in the most important applications of the G -matrix E lies below the Fermi level. For instance, in the Brückner Hartree-Fock theory (see Chap. 5) we must calculate $G_{ab,ij}^{E=\epsilon_i+\epsilon_j}$ with $i, j < \epsilon_F$. We will therefore treat the hard core for this case in a very simplified but explicitly solvable model which, however, shows the essential features.

Let us consider a large nucleus, the interior of which will be presumed to be not very much different from the situation where we consider an infinitely large nucleus, usually termed nuclear matter. In addition, we will make a further crude assumption, namely, that the interaction consists of a hard core only:

$$V(r) = 0 \text{ for } r \geq c \quad \text{and} \quad V(r) = \infty \text{ for } r < c.$$

Furthermore, we assume that we have two distinguishable particles, for example, a proton and a neutron, which will be sufficient as a demonstration of the principle. Let us now write Eq. (4.43) in the form:

$$(E - H_0)|\psi_{ab}\rangle = (E - H_0 + Q_F G)|ab\rangle, \quad (4.49)$$

which in our case ($E = \epsilon_i + \epsilon_j$) specializes to

$$(\epsilon_i + \epsilon_j - H_0)|\psi_{ij}\rangle = Q_F G |ij\rangle = Q_F \bar{v} |\psi_{ij}\rangle, \quad (4.50)$$

where for the last equality use has been made of (4.43). The general solution of (4.50) is much harder than the solution of an ordinary Schrödinger equation because of the operator Q_F , which is a nonlocal integral operator. This is best seen in the \mathbf{r} -representation

$$\langle \mathbf{r}_1 \mathbf{r}_2 | Q_F | \mathbf{r}'_1 \mathbf{r}'_2 \rangle = \sum_{\substack{m < n \\ > \epsilon_F}} \langle \mathbf{r}_1 \mathbf{r}_2 | mn \rangle \langle mn | \mathbf{r}'_1 \mathbf{r}'_2 \rangle. \quad (4.51)$$

In nuclear matter it is clear that the single-particle wave functions that appear in (4.50) are plane waves. For simplicity we will also assume that the single-particle energies are purely kinetic energies (in a surrounding medium this will not gener-

ally be the case). Because of translational invariance, for nuclear matter the center of mass motion is trivial, and we have:

$$\langle \mathbf{r}_1 \mathbf{r}_2 | \psi_{ij} \rangle \rightarrow \psi_{\mathbf{p}_1 \mathbf{p}_2}(\mathbf{r}_1 \mathbf{r}_2) = \frac{1}{(2\pi)^{3/2}} e^{i\mathbf{P}\mathbf{R}} \psi_{\mathbf{p}_k}(\mathbf{r}). \quad (4.52)$$

Here the following transformations to relative and center of mass coordinates have been made:

$$\begin{aligned} \mathbf{P} &= \mathbf{p}_1 + \mathbf{p}_2, & \mathbf{k} &= \frac{1}{2}(\mathbf{p}_1 - \mathbf{p}_2), \\ \mathbf{R} &= \frac{1}{2}(\mathbf{r}_1 + \mathbf{r}_2), & \mathbf{r} &= \mathbf{r}_1 - \mathbf{r}_2. \end{aligned} \quad (4.53)$$

With all these assumptions, Eq. (4.50) acquires the form

$$\begin{aligned} \frac{\hbar^2}{2m} \{ p_1^2 + p_2^2 + \Delta_1 + \Delta_2 \} \psi_{\mathbf{p}_1 \mathbf{p}_2}^*(\mathbf{r}_1 \mathbf{r}_2) \\ = \frac{1}{(2\pi)^6} \int_{p, p' > k_F} d^3 p d^3 p' \int d^3 r'_1 d^3 r'_2 e^{i(\mathbf{p}\mathbf{r}_1 + \mathbf{p}'\mathbf{r}_2 - \mathbf{p}'\mathbf{r}'_1 - \mathbf{p}'\mathbf{r}'_2)} v(|\mathbf{r}'_1 - \mathbf{r}'_2|) \psi_{\mathbf{p}_1 \mathbf{p}_2}(\mathbf{r}'_1 \mathbf{r}'_2). \end{aligned} \quad (4.54)$$

Transforming this equation to relative and center of mass coordinates according to (4.53) we obtain

$$\begin{aligned} \frac{\hbar^2}{m} \{ k^2 + \Delta_r \} e^{i\mathbf{P}\mathbf{R}} \psi_{\mathbf{p}_k}(\mathbf{r}) \\ = \frac{1}{(2\pi)^6} \int \int_{\left| \frac{\mathbf{P}}{2} \pm \mathbf{p} \right| > k_F} d^3 \mathcal{P} d^3 p \int \int d^3 R' d^3 r' e^{i[\mathcal{P}\mathbf{R} + \mathbf{p}(\mathbf{r} - \mathbf{r}')] - i\mathcal{P}\mathbf{R}'} v(r') e^{i\mathbf{P}\mathbf{R}'} \psi_{\mathbf{p}_k}(\mathbf{r}'). \end{aligned} \quad (4.55)$$

Since $\mathcal{P}^2 = p^2 + 2pp' \cos \theta + p'^2$, we see that even under the restriction $p, p' > k_F$, \mathcal{P} can take on all values from 0 to ∞ . The integration over \mathbf{R}' gives $\delta(\mathbf{P} - \mathcal{P})$ and we can therefore also perform the \mathcal{P} integral. We are thus left with the following equation.

$$\frac{\hbar^2}{m} \{ k^2 + \Delta_r \} \psi_{\mathbf{p}_k}(\mathbf{r}) = \frac{1}{(2\pi)^3} \int_{\left| \frac{\mathbf{P}}{2} \pm \mathbf{p} \right| > k_F} d^3 p e^{i\mathbf{p}\mathbf{r}'} \int d^3 r' e^{-i\mathbf{p}\mathbf{r}'} v(r') \psi_{\mathbf{p}_k}(\mathbf{r}'). \quad (4.56)$$

This equation is not only more complicated than a usual two-particle equation because of its integrodifferential structure, but also because the wave function has a nontrivial dependence on the center of mass momentum \mathbf{P} . For our purpose, it will be sufficient to evaluate it at $\mathbf{P} = 0$. Furthermore, we can decompose Eq. (4.56) into partial waves [Me 61 Chap. X, Sec. 8]:

$$\psi_{\mathbf{k}}(\mathbf{r}) = \sum_{lm} \psi_{\mathbf{k}l}(r) Y_{lm}(\theta, \varphi). \quad (4.57)$$

Considering the equation for the s -wave and splitting the integral in (4.56) into two parts, $\int_{k_F}^{\infty} = \int_0^{\infty} - \int_0^{k_F}$, we obtain:

$$\begin{aligned} \frac{\hbar^2}{m} \left(k^2 + \frac{\partial^2}{\partial r^2} + \frac{2}{r} \frac{\partial}{\partial r} \right) \psi_{\mathbf{k}}(r) \\ = v(r) \psi_{\mathbf{k}}(r) - (4\pi)^2 \int_0^{k_F} \frac{dp p^2}{(2\pi)^3} j_0(pr) \int_0^{\infty} dr' r'^2 j_0(pr') v(r') \psi_{\mathbf{k}}(r'), \end{aligned} \quad (4.58)$$

where j_0 is the lowest-order spherical Bessel function. With $\psi_{\mathbf{k}}(r) = (1/r)u_k(r)$, we finally get:

$$\frac{\hbar^2}{m} \left(k^2 + \frac{\partial^2}{\partial r^2} \right) u_k(r) = v(r)u_k(r) - k_F \int_0^\infty dr' \chi(r', r)v(r')u_k(r'), \quad (4.59)$$

with

$$\chi(r, r') = \frac{2rr'}{\pi k_F} \int_0^{k_F} dp p^2 j_0(pr)j_0(pr') = \frac{1}{\pi \cdot k_F} \left\{ \frac{\sin k_F(r-r')}{r-r'} - \frac{\sin k_F(r+r')}{r+r'} \right\}. \quad (4.60)$$

It is now convenient to introduce the following dimensionless quantities:

$$x = k_F \cdot r; \quad x' = k_F \cdot r'; \quad K = \frac{k}{k_F}; \quad \frac{mv}{k_F^2 \hbar^2} = w. \quad (4.61)$$

This leads to

$$\left(\frac{d^2}{dx^2} + K^2 \right) u_K(x) = w(x)u_K(x) - \int_0^\infty dx' \chi(x, x')w(x')u_K(x'). \quad (4.62)$$

For a square well barrier of finite height, the wave function and its first derivative are continuous at the edge of the barrier. We can be easily convinced that the first derivative of the wave function becomes discontinuous at the edge as the barrier height goes to infinity, that is:

$$u'_K(c + \epsilon) = A, \quad (4.63)$$

where A is a constant, $r = c$ is the radius of the hard core, and ϵ is an infinitesimally small positive quantity. In order to produce such a discontinuity for u of Eq. (4.62), the product $w \cdot u$ must be proportional to a δ function for $r = c$. Since for $r > c$ the potential w is zero and u finite, the product $w \cdot u$ vanishes outside the hard core radius. The wave function u cannot penetrate inside the infinite hard core ($u=0$); since there $w=\infty$, the product $w \cdot u$ may be finite and we can write, with $c' = ck_F$:

$$w(x)u_K(x) = A\delta(x - c') + l(x)\theta(c' - x). \quad (4.64)$$

The function $l(x)$ can be determined from the requirement that for $x < c'$ the left hand side of Eq. (4.62) must be zero, since $u \equiv 0$ for $x < c'$. We therefore have from Eq. (4.62) with Eq. (4.64)

$$l(x) = A\chi(x, c') + \int_0^{c'} dx' \chi(x, x')l(x') \quad \text{for } x < c'. \quad (4.65)$$

Since the hard core is usually rather small ($c' = 0.57$ at the nuclear matter density for $c = 0.4$ fm), we can develop the kernel in Eq. (4.65) and obtain from Eq. (4.60)

$$\chi(x, x') \rightarrow \frac{2xx'}{3\pi} \quad x < c'; \quad x' < c'. \quad (4.66)$$

With Eqs. (4.66) and (4.65), we see that $l(x)$ is of order c'^2 , whereas any integral over $l(x)$ will be of order c'^3 . In (4.62) we can therefore neglect the second term in (4.64) to obtain a result which is correct to first order in c' . We obtain

$$\begin{aligned} \left(\frac{d^2}{dx^2} + K^2 \right) u_K(x) &\simeq A[\delta(x - c') - \chi(x, c')] + A\chi(x, c')\theta(c' - x) \\ &= A \frac{2xc'}{\pi} \int_1^\infty dp p^2 j_0(px)j_0(pc') + A\chi(x, c')\theta(c' - x) \\ &= F(x), \end{aligned} \quad (4.67)$$

where we have made use of the identity

$$\int_0^\infty dp p^2 j_l(pr) j_l(pr') = \frac{\pi}{2r^2} \delta(r - r'). \quad (4.68)$$

We remark that since the r.h.s. of Eq. (4.67) is only correct to first order in c' , it is somewhat arbitrary whether to include the second term of the r.h.s. of Eq. (4.67), which is of second order. The general solution of Eq. (4.67),

$$u_K(x) = \frac{\sin Kx}{K} \int_{c'}^x dy F(y) \cos Ky - \frac{\cos Kx}{K} \int_{c'}^x dy F(y) \sin Ky, \quad (4.69)$$

is obviously zero for $x = c'$. Equation (4.69) is the solution to Eq. (4.67), as can be easily verified by direct insertion.

The only unknown in the solution (4.67, 4.69) is the constant A , which we are going to determine by the requirement (4.46) that u_K has to approach asymptotically the unperturbed value. To this end we will first show that the second integral on the r.h.s. of Eq. (4.69) is zero in the limit $x \rightarrow \infty$. Since we considered our solution in the limit of very small c' , we can take the integral from 0 to ∞ instead of from c' to ∞ , and thus have

$$\begin{aligned} \frac{1}{K} \int_0^\infty dy F(y) \sin Ky &= \frac{2Ac'}{\pi} \int_1^\infty dp p^2 j_0(p c') \int_0^\infty dy y^2 j_0(Ky) j_0(py) \\ &= \frac{2Ac'}{\pi} \int_1^\infty dp p^2 j_0(p c') \frac{\pi}{2Kp} \delta(K - p) \\ &= 0. \end{aligned} \quad (4.70)$$

The last integral vanishes because p is outside the Fermi sphere, whereas K is inside. For $r \rightarrow \infty$ we therefore find the result that the wave function approaches the unperturbed result ($\psi \rightarrow j_0$) [see Eq. (4.46)] only if

$$\int_0^\infty dy F(y) \cos Ky = 1, \quad (4.71)$$

and we therefore get, using Eq. (4.67):

$$A = \left\{ \cos Kc' - \int_0^\infty dy \cos(Ky) \chi(y, c') \right\}^{-1}, \quad (4.72)$$

which completes the determination of the wave function $u_K(x)$ of Eqs. (4.67) and (4.69) for a hard sphere potential.

In Fig. 4.3 we show the solution $\psi_k(r)$ for $k=0$, which reveals several interesting features. The wave function vanishes inside the hard core. With (4.69) and (4.72) it can easily be seen that it approaches rapidly (by

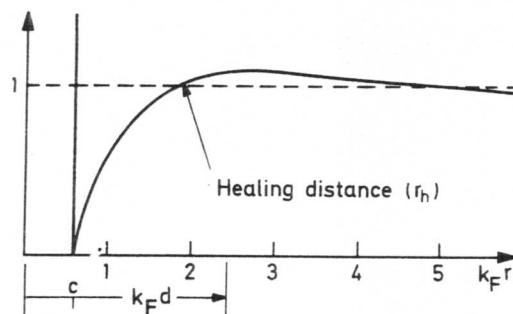


Figure 4.3. S-wave solution of the Bethe-Goldstone equation for a pure hard core potential in nuclear matter. For details of the figure, see text. (From [FW 71].)

damped oscillations) the value of the wave function for the noninteracting pair (which is itself equal to 1 for $k=0$ and all values of r). The r value at which the wave function first attains the unperturbed value is called the *healing distance* r_h , because it is only for distances smaller than this value that the hard core wave function differs appreciably from the unperturbed wave function. For the values chosen in our example, the healing distance is:

$$k_F r_h \approx 1.9. \quad (4.73)$$

(This value is more or less independent of \mathbf{P} and k , as a more general study of Eq. (4.59) shows [WG 67a]). It is important to note that this value is appreciably smaller than the average interparticle distance d in nuclear matter defined by the expression $1/d^3 \equiv N/V = (4 \int_{p < \hbar k_F} d\mathbf{p} \int d\mathbf{r})/V(2\pi\hbar)^3 = 2k_F^3/3\pi^2$, which would yield for the interparticle distance:

$$k_F d = \left(\frac{3\pi^2}{2} \right)^{1/3} = 2.46. \quad (4.74)$$

On the average, therefore, the nucleons return after a collision to their independent particle wave function *before* the next collision takes place. A great portion of the nuclear wave function therefore consists of a determinant built out of independent particle wave functions. This can be considered as a justification of the independent particle model—the Pauli principle suppressing the low momentum components of the scattering process, that is, suppressing the long range correlations. The nucleons thus move through the nucleus most of the time as independent particles because of the exclusion principle. That this can happen in a strongly interacting but dilute Fermi gas was first pointed out by Weisskopf [We 50].

We have seen that the solution of the Bethe–Goldstone Equation (4.39) is far from trivial, and one can easily imagine that the task can become tremendously difficult for finite nuclei, where the wave functions are no longer plane waves and translational invariance is also lost. Several approximation schemes have therefore been currently applied. As we have seen in our example, it is the projection operator Q_F which makes things so difficult. It has been proposed, that the Pauli principle should be treated only perturbatively in the *separation method* of Moszkowski and Scott [MS 60] and in the *reference spectrum method* of Bethe, Brandow, and Petschek [BBP 63]. We do not wish to go into details of these approximation methods and refer the reader to the text books of Brown [Br 64], de Shalit and Feshbach [SF 74], and the review article by Bethe [Be 71], in which these and further methods, like the use of the oscillator basis for the solution of the Bethe–Goldstone equation in finite nuclei, are explained in detail. Here we want to mention only one other approximation scheme which has recently become very successful in connection with the Brückner–Hartree–Fock theory (see Chap. 5). This is the *local density approximation* by Negele [Ne 70, 75], originally introduced by Brückner, Gammel, and Weitzner [BGW 58] and Brückner, Lockett, and Rotenberg [BLR 61]. The assumption is that the G -matrix at any place in a finite nucleus is the same as that for nuclear matter at the same density, so that locally one can