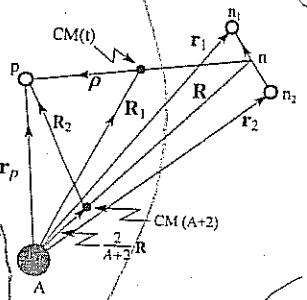
Appendix 7.A ZPF and Pauli principle at the basis of medium polarization effects: self–energy, vertex corrections and induced interaction

In keeping with a central objective of the formulation of quantum mechanics, namely that the basic concepts on which it is based relate directly to experiment (Heisenberg), elementary modes of nuclear excitation (single–particle, collective vibrations and rotations), are solidly anchored on observation (inelastic and Coulomb excitation, one– and two–particle transfer reactions). Of all quantal phenomena, zero point fluctuations (ZPF), closely connected with virtual states, are likely to be most representative of the essential difference existing between quantum and classical mechanics. In fact, ZPF are intimately connected with the complementary principle (Bohr), and thus with the indeterminacy (Heisenberg) and non–commutative (Born, Jordan) relations, and with the probabilistic interpretation (Born) of the (modulus squared) of the wavefunctions, solution of Schrödinger's or Dirac's equations.

Pauli principle brings about essential modifications of the virtual fluctuations of the many–body system, modifications which are instrumental in the dressing and interweaving of the elementary modes of excitation (see Figs. 7.A.1 and 7.A.2); within the present context, see also Schrieffer (1964).

Appendix 7.B Coherence and effective formfactors



$r = r_1 - r_2$

$r = r_1 - r_2$ (relative distance between the neutrons)

$R = \frac{r_1 + r_2}{2}$ (coord. of the CM of the dineutron)

$\rho = r_p - \frac{r_1 + r_2}{2} = r_p - R$ (distance between the CM of the dineutron and the proton)

$R_2 = r_p - \frac{r_1 + r_2}{A+2}$ (distance of the proton from the CM of the system A+2)

$R_1 = \frac{r_1 + r_2 + r_3}{3}$ (coord. of the CM of the triton)

Figure 7.B.1:

Again we assume that the reaction is direct, and that it is adequately described by first–order distorted–wave Born approximation.

To be specific, we will concentrate on (t, p) reaction, namely reactions of the type $A(\alpha, \beta)B$ where $\alpha = \beta + 2$ and $B = A + 2$.

The intrinsic wave functions are in this case

$$\begin{aligned}\psi_\alpha =&\psi^{J_i}_{M_i}(\xi_A)\sum_{ss'_f}\left[\chi^s(\sigma_\alpha)\chi^{s'_f}(\sigma_\beta)\right]^{s_i}_{M_{s_i}}\phi^{L=0}_t(\sum_{i<j}|\vec{r}_i-\vec{r}_j|)\\ =&\psi^{J_i}_{M_i}(\xi_A)\sum_{M_s M'_{s_f}}(sM'_{s_f}s'_f M'_{s_f}|s_i M_{s_i})\chi^s_{M_s}(\sigma_\alpha)\chi^{s_f}_{M_{s_f}}(\sigma_\beta)\\ &\times\phi^{L=0}_t(\sum_{i<j}|\vec{r}_i-\vec{r}_j|)\end{aligned}\tag{7.B.1}$$

while

$$\begin{aligned}\psi_\beta =&\psi^{J_f}_{M_f}(\xi_{A+2})\chi^{s_f}_{M_{s_f}}(\sigma_\beta)\\ =&\sum_{\substack{n_1l_1j_1\\ n_2l_2j_2}}B(n_1l_1j_1,n_2l_2j_2);JJ'_iJ_f)\left[\phi^J(j_1j_2)\phi^{J'_i}(\xi_A)\right]^{J_f}_{M_f}\\ &\times\chi^{s_f}_{M_{s_f}}(\sigma_\beta)\end{aligned}\tag{7.B.2}$$

But from eq. (7.B.2) is easy to see that the spectroscopic amplitude B is equal to

$$\begin{aligned}&B(n_1l_1j_1,n_2l_2j_2);JJ'_iJ_f)\\ &=\left\langle\psi^{J_f}(\xi_{A+2})\left|\left[\phi^J(j_1j_2)\phi^{J_i}(\xi_A)\right]^{J_f}\right.\right\rangle\end{aligned}\tag{7.B.3}$$

where

$$\phi^J(j_1j_2)=\frac{\left[\phi_{j_1}(\vec{r}_1)\phi_{j_2}(\vec{r}_2)\right]^J-\left[\phi_{j_1}(\vec{r}_2)\phi_{j_2}(\vec{r}_1)\right]^J}{\sqrt{1+\delta(j_1,j_2)}}\tag{7.B.4}$$

is an antisymetrized, normalized wave function of the two transferred particles. The function $\chi^s_{M_s}(\sigma_\beta)$ appearing both in eq. (7.B.1) and (7.B.2) is the spin wave function of the proton while $\chi^s(\sigma_\alpha)$ is equal to

$$\chi^s(\sigma_\alpha)=[\chi^{s_1}(\sigma_{n_1})\chi^{s_2}(\sigma_{n_2})]^s\tag{7.B.5}$$

is the spin function of the two–neutron system.

The function $\phi^{L=0}_t$ describes the internal degree of freedom of the triton. A good description of this system is obtained by using a wave function symmetric in the coordinates of all particles, i.e.

$$\begin{aligned}\phi^{L=0}_t(\sum_{i<j}|\vec{r}_i-\vec{r}_j|)&=N_t\,e^{[(r_1-r_2)^2+(r_1-r_p)^2+(r_2-r_p)^2]}\\ &=\phi_{000}(\vec{r})\phi_{000}(\vec{\rho})\\ \phi_{000}(\vec{r})&=R_{nl}(\nu^{1/2}r)Y_{lm}(\hat{r})\end{aligned}\tag{7.B.6}$$

The coordinate $\vec{\rho}$ is the radius vector which measures the distance between the center of mass of the dineutron and the proton, while the vector $\vec{r}$ is the dineutron relative coordinate

~~$\vec{r}=\vec{r}_1-\vec{r}_2$ (relative distance between the neutrons)~~ (7.B.7a)

$$\vec{R} = \frac{\vec{r}_1 + \vec{r}_2}{2} \quad \text{(coord. of the CM of the dineutron)} \qquad \text{(7.B.7b)}$$

$$\vec{\rho} = \vec{r}_p - \frac{\vec{r}_1 + \vec{r}_2}{2} \quad \text{(distance between the CM of the dineutron and the proton)} \qquad \text{(7.B.7c)}$$

$$\vec{R}_2 = \vec{r}_p - \frac{\vec{r}_1 + \vec{r}_2}{A+2} \quad \text{(distance of the proton from the CM of the system A+2)} \qquad \text{(7.B.7d)}$$

$$\vec{R}_1 = \frac{\vec{r}_p + \vec{r}_1 + \vec{r}_2}{3} \quad \text{(coord. of the CM of the triton)} \qquad \text{(7.B.7e)}$$

To obtain the DWBA cross section we have to calculate the integral

$$T = \int d\xi_A \, d\vec{r}_1 \, d\vec{r}_2 \, d\vec{r}_p \chi_p^{(-)}(\vec{R}_2)\psi_\beta^*(\xi_{A+2}, \sigma_\beta)V'_\beta\psi_\alpha(\xi_A, \sigma_\alpha, \sigma_\beta)\psi_t^{(+)}(\vec{R}_1) \qquad \text{(7.B.8)}$$

Instead of integrating over $\xi_A, \vec{r}_1, \vec{r}_2$ and $\vec{r}_p$ we would integrate over $\xi_A, \vec{r}, \vec{r}$ and $\vec{r}_p$. The Jacobian of the transformation is equal to 1, i.e. $\partial(\vec{r}_1, \vec{r}_2)/\partial(\vec{r}, \vec{r}) = 1$.

To carry out the integral (7.B.8) we transform the wave function (7.B.4) into center of mass and relative coordinates. If we assume that both $\phi_{j_1}(\vec{r}_1)$ and $\phi_{j_2}(\vec{r}_2)$ are harmonic oscillator wave functions, this transformation can easily carried with the aid of tha Moshinsky brackets. If $|n_1 l_1, n_2 l_2; \lambda\mu\rangle$ is acomplete system of wave functions in the harmonic oscillator basis, depending on $\vec{r}_1$ and $\vec{r}_2$ and $|nl, NL; \lambda\mu\rangle$ is the corresponding one depending on $\vec{r}$ and $\vec{R}$, we can write

$$\begin{aligned} |n_1 l_1, n_2 l_2; \lambda\mu\rangle &= \sum_{nlNL} |nl, NL; \lambda\mu\rangle\langle nl, NL; \lambda\mu|n_1 l_1, n_2 l_2; \lambda\mu\rangle \\ &= \sum_{nlNL} |nl, NL; \lambda\mu\rangle\langle nl, NL; \lambda\mu|n_1 l_1, n_2 l_2; \lambda\rangle \end{aligned} \qquad \text{(7.B.9)}$$

The labels n, l are the principal and angular momentum quantum numbers of the relative motion, ehile N, L are the corresponding ones corresponding to the center of mass motion of the two–neutron system. Because of energy and parity conservation we have

$$\begin{aligned} 2n_1 + l_1 + 2n_2 + l_2 &= 2n + l + 2N + L \\ (-1)^{l_1+l_2} &= (-1)^{l+L} \end{aligned} \qquad \text{(7.B.10)}$$

The coefficients $\langle nl, NL, L|n_1 l_1, n_2 l_2, L\rangle$ are tabulated and were first discussed by M. Moshinsky in Nucl. Physics, 13 (1959) 104.

With the help of eq.(7.B.9) we can write the wave function $\psi^{J_f}_{M_f}(\xi_{A+2})$ as

$$
\begin{aligned}
\psi^{J_f}_{M_f}(\xi_{A+2}) &= \sum_{\substack{n_1 l_1 j_1 \\ n_2 l_2 j_2 \\ JJ_i}} B(n_1 l_1 j_1, n_2 l_2 j_2; JJ_i' J_f) \left[\phi^{j}(j_1 j_2)\phi^{J_i'}(\xi_A)\right]^{J_f}_{M_f} \\
&= \sum_{\substack{n_1 l_1 j_1 \\ n_2 l_2 j_2}} \sum_{JJ_i} B(n_1 l_1 j_1, n_2 l_2 j_2; JJ_i' J_f) \\
&\times \sum_{M_J M_{J_i}'} \langle J M_J J_i' M_{J_i} | J_f M_{J_f}\rangle \psi^{J_i'}_{M_{J_i}'}(\xi_A) \\
&\times \sum_{LS'} \langle S' L J | j_1 j_2 J\rangle \sum_{M_L M_S'} \langle L M_L S' M_S' | J M_J\rangle \chi^{S'}_{M_S'}(\sigma_\alpha) \\
&\times \sum_{nl N\Lambda} \langle nl, N\Lambda, L | n_1 l_1, n_2 l_2, L\rangle \\
&\times \sum_{m_l M_\Lambda} \langle l m_l \Lambda M_\Lambda | L M_L\rangle \phi_{nlm_l}(\vec{r})\phi_{N\Lambda M_\Lambda}(\vec{R})
\end{aligned}
\tag{7.B.11}
$$

Integration over $\vec{r}$ gives

$$
\langle \phi_{nlm_l}(\vec{r}) | \phi_{000}(\vec{r})\rangle = \delta(l,0)\delta(m_l,0)\Omega_n \tag{7.B.12}
$$

where

$$
\Omega_n = \int R_{nl}(\nu_1^{1/2} r) R_{00}(\nu_2^{1/2} r) r^2\, dr \tag{7.B.13}
$$

Note that there is no selection rule in the principal quantum number n, as the potential in which the two neutrons move in the triton has a frequency ν_2 which is different from the one that the two neutrons are subjected to, when moving in the system A.

Integration over ξ_A and multiplication of the spin functions gives

$$
\begin{aligned}
(\psi^{J_i}_{M_{J_i}}, (V(\vec{r}_1) + V(\vec{r}_2) + V(\vec{r}_p) - U)\psi^{J_i'}_{M_{J_i}'}) &= \delta(J_i, J_i')\delta(M_{J_i}, M_{J_i'}) V_{eff}(\vec{\rho}) \\
(\chi^{S}_{M_S}(\sigma_\alpha), \chi^{S'}_{M_{S'}}(\sigma_\alpha)) &= \delta(S, S')\delta(M_S, M_{S'}) \\
(\chi^{S_f}_{M_{S_f}}(\sigma_\beta), \chi^{S_f'}_{M_{S_f'}}(\sigma_\beta)) &= \delta(S_f, S_f')\delta(M_{S_f}, M_{S_f'})
\end{aligned}
\tag{7.B.14}
$$

The integral (7.B.8) can now be written as

$$
\begin{aligned}
T = &\sum_{\substack{n_1 l_1 j_1 \\ n_2 l_2 j_2}} \sum_{JM_J} \sum_{nN} \sum_{S} B(n_1 l_1 j_1, n_2 l_2 j_2; JJ_i' J_f) \\
&\times \langle J M_J J_i M_{J_i} | J_f M_{J_f}\rangle \langle S L J | j_1 j_2 J\rangle \\
&\times \langle L M_L S M_S | J M_J\rangle \langle n0, NL, L | n_1 l_1, n_2 l_2, L\rangle \\
&\times \langle S M_S S_f M_{S_f} | S_i M_{S_i}\rangle \Omega_n \\
&\times \int d\vec{R}\, d\vec{r}_p \chi_t^{(+)*}(\vec{R}_1)\phi^*_{NLM_L}(\vec{R}) V_{eff}(\vec{\rho})\phi_{000}(\vec{\rho})\chi_t^{(+)}(\vec{R}_1)
\end{aligned}
\tag{7.B.15}
$$

where we have approximated V_β' by an effective interaction V_{eff} depending only on $\rho = |\vec{\rho}|$. It is important to point out that the two–body interaction would act on the

two–particle system at once, but the single particle potential would act on each particle independently. The reason why we can neglect the successive transfer of the nucleons (two–step process) is because the two neutrons in the triton are very strongly correlated and they build to a large extent a unity.

We now define the two–nucleon transfer form factor as

$$u_{LSJ}^{j_i J_f}(R) = \sum_{n_1 l_1 j_1} B(n_1 l_1 j_1, n_2 l_2 j_2; J J_i J_f) \langle S L J | j_1 j_2 J \rangle \langle n0, NL, L | n_1 l_1, n_2 l_2; L \rangle \Omega_n R_{nL}(R) \tag{7.B.16}$$

We can now rewrite eq. (7.B.15) as

$$T = \sum_J \sum_L \sum_S (J M_J J_i M_{J_i} | J_f M_{J_f}) \langle S M_S S_f M_{S_f} | S_i M_{S_i} \rangle (L M_L S M_S | J M_J) \times \int d\vec{R}\, d\vec{r}_p \chi_p^{*(-)}(\vec{R}_2) u_{LSJ}^{j_i J_f}(R) Y_{LM_L}^* V(\rho) \phi_{000}(\vec{\rho}) \chi_t^{(+)}(\vec{R}_1) \tag{7.B.17}$$

Because the di–neutron has $S = 0$, we have that

$$(L M_L 00 | J M_J) = \delta(J, L) \delta(M_L, M_J) \tag{7.B.18}$$

and the summations over S and L disappear from eq. (7.B.17).

The integral to be carried out in eq. (7.B.17) is six–dimensional, and is a formidable task to calculate it exactly (one of these integrals takes $\approx$ 5 hs in a CDC 6600 computer).

One can make also here the zero range approximation, i.e.

$$V(\rho) \phi_{000}(\vec{\rho}) = D_0 \delta(\vec{\rho}) \tag{7.B.19}$$

This means that the proton interacts with the center of mass of the di–neutron, only when they are at the same point in space.

From eqs. (??) we obtain

$$\begin{aligned} \vec{R} &= \vec{R}_1 = \vec{r} \\ \vec{R}_2 &= \frac{A}{A+2} \vec{R} \end{aligned} \tag{7.B.20}$$

Then eq. (7.B.15) can be written as

$$T = D_0 \sum_L (L M_L J_i M_{J_i} | J_f M_{J_f}) \times \int d\vec{R} \chi_p^{*(-)}(\frac{A}{A+2}\vec{R}) u_L^{j_i J_f}(R) Y_{LM_L}^*(\hat{R}) \chi_t^{(+)}(\vec{R}) \tag{7.B.21}$$

From eq. (7.B.21) it is seen that the change in parity implied by the reaction is given by $\Delta\Pi = (-1)^L$. Consequently, the selection rules for (t, p) and (p, t) reactions are

$$\begin{aligned} \Delta S &= 0 \\ \Delta J &= \Delta L = L \\ \Delta\Pi &= (-1)^L \end{aligned} \tag{7.B.22}$$

i.e. only normal parity states are excited.

The integral appearing in eq. (7.B.21) has the same structure as the DWBA integral (??) (Fourth Lecture) which was derived in the case of one–nucleon transfer reactions.

The difference between the two processes manifest itself through the different structure of the two form factors. While $u_l(r)$ appearing in equation (??) is a single–particle bound state wave function, $u_L^{J_i J_f}$ is a coherent summation over the center of mass states of motion of the two transferred neutrons.

Appendix 7.C Relative importance of successive and simultaneous transfer and non-orthogonality corrections

In what follows we discuss the relative importance of successive and simultaneous two-neutron transfer and of non-orthogonality corrections associated with the reaction

$$\alpha \equiv a(=b+2)+A \rightarrow b+B(=A+2) \equiv \beta \tag{7.C.1}$$

in the limits of independent particles and of strongly correlated Cooper pairs, making use for simplicity of the semiclassical approximation (for details cf. Broglia and Winther (2004),? and refs. therein), in which case the two-particle transfer differential cross section can be written as

$$\frac{d\sigma_{\alpha\rightarrow\beta}}{d\Omega} = P_{\alpha\rightarrow\beta}(t=+\infty)\sqrt{\left(\frac{d\sigma_\alpha}{d\Omega}\right)_{el}}\sqrt{\left(\frac{d\sigma_\beta}{d\Omega}\right)_{el}}, \tag{7.C.2}$$

where P is the absolute value squared of a quantum mechanical transition amplitude. It gives the probability that the system at $t=+\infty$ is found in the final channel. The quantities $(d\sigma/d\Omega)_{el}$ are the classical elastic cross sections in the center of mass system, calculated in terms of the deflection function, namely the functional relating the impact parameter and the scattering angle.

The transfer amplitude can be written as

$$a(t=+\infty) = a^{(1)}(\infty) - a^{(NO)}(\infty) + \bar{a}^{(2)}(\infty), \tag{7.C.3}$$

where $\bar{a}^{(2)}(\infty)$ at $t=+\infty$ labels the successive transfer amplitude expressed in the post-prior representation (see below). The simultaneous transfer amplitude is given by (see Fig. A.1(I))

$$a^{(1)}(\infty) = \frac{1}{i\hbar}\int_{-\infty}^{\infty} dt(\psi^b\psi^B, (V_{bB} - < V_{bB} >)\psi^a\psi^A) \times \exp[\frac{i}{\hbar}(E^{bB} - E^{aA})t]$$
$$\approx \frac{2}{i\hbar}\int_{-\infty}^{\infty} dt\left(\phi^{B(A)}(S^B_{(2n)}; \vec{r}_{1A}, \vec{r}_{2A}), U(r_{1b})e^{i(\sigma_1+\sigma_2)}\phi^{a(b)}(S^a_{(2n)}; \vec{r}_{1b}, \vec{r}_{2b})\right)\exp[\frac{i}{\hbar}(E^{bB} - E^{aA})t + \gamma(t)] \tag{7.C.4}$$

where

$$\sigma_1 + \sigma_2 = \frac{1}{\hbar}\frac{m_n}{m_A}(m_{aA}\vec{v}_{aA}(t) + m_{bB}v_{bB}(t)) \cdot (\vec{r}_{1\alpha} + \vec{r}_{2\alpha}), \tag{7.C.5}$$

in keeping with the fact that $\exp(i(\sigma_1+\sigma_2))$ takes care of recoil effects (Galilean transformation associated with the mismatch between entrance and exit channels).

The phase $\gamma(t)$ is related with the effective $Q-$value of the reaction. In the above expression, ϕ indicates an antisymmetrized, correlated two-particle (Cooper pair) wavefunction, $S(2n)$ being the two-neutron separation energy (see Fig. A1), $U(r_{1b})$ being the single particle potential generated by nucleus b ($U(r) = \int d^3r' \rho^b(r') v(|r - r'|)$). The contribution arising from non-orthogonality effects can be written as (see Fig. A1(II))

$$a^{(NO)}(\infty) = \frac{1}{i\hbar}\int_{-\infty}^{\infty} dt(\psi^b\psi^B, (V_{bB} - < V_{bB} >)\psi^f\psi^F)(\psi^f\psi^F, \psi^a\psi^A)\exp[\frac{i}{\hbar}(E^{bB} - E^{aA})t]$$

$$\approx \frac{2}{i\hbar}\int_{-\infty}^{\infty} \phi^{B(F)}(S^B_{(n)}, \vec{r}_{1A}), U(r_{1b})e^{i\sigma_1}(\phi^{f(b)}(S^f(n), \vec{r}_{1b})\phi^{F(A)}(S^F(n), \vec{r}_{2A})e^{i\sigma_2}\phi^{a(f)}(S^a(n), \vec{r}_{2b}))\exp[\frac{i}{\hbar}(E^{bB} - E^{aA})t + \gamma(t)] \quad \text{(7.C.6)}$$

the reaction channel $f = (b + 1) + F(= A + 1)$ having been introduced, the quantity $S(n)$ being the one-neutron spearation energy (see Fig. A1). The summation over $f(\equiv a'_1, a'_2)$ and $F(\equiv a_1, a_2)$ involves a restricted number of states, namely the valence shells in nuclei B and a.

The successive transfer amplitude $\bar{a}^{(2)}_{\infty}$ written making use of the post-prior representation is equal to (see Fig. A1(III))

$$\bar{a}^{(2)}(\infty) = \frac{1}{i\hbar}\int_{-\infty}^{\infty} dt(\psi^b\psi^B, (V_{bB} - < V_{bB} >)e^{i\sigma_1}\psi^f\psi^F) \times \exp[\frac{i}{\hbar}(E^{bB} - E^{fF})t + \gamma_1(t)]$$

$$\times\frac{1}{i\hbar}\int_{-\infty}^{t} dt'(\psi^f\psi^F, (V_{fF} - < V_{fF} >)e^{i\sigma_2}\psi^a\psi^A > \times\exp[\frac{i}{\hbar}(E^{fF} - E^{aA})t' + \gamma_2(t)] \quad \text{(7.C.7)}$$

To gain insight into the relative importance of the three terms contributing to Eq. 7.C.3 we discuss two situations, namely, the independent-particle model and the strong-correlation limits.

7.C.1 Independent particle limit

In the independent particle limit, the two transferred particles do not interact among themselves but for antisymmetrization. Thus, the separation energies fulfill the relations (see Fig. A2)

$$S^B(2n) = 2S^B(n) = 2S^F(n), \quad \text{(7.C.8)}$$

and

$$S^a(2n) = 2S^a(n) = 2S^f(n). \quad \text{(7.C.9)}$$

In this case

$$\phi^{B(A)}(S^B(2n), \vec{r}_{1A}, \vec{r}_{2A}) = \sum_{a_1a_2} \phi^{B(F)}_{a_1}(S^B(n), \vec{r}_{1A})\phi^{F(A)}_{a_2}(S^F(n), \vec{r}_{2a}), \quad \text{(7.C.10)}$$

and

$$\phi^{a(b)}(S^a(2n), \vec{r}_{1b}, \vec{r}_{2b}) = \sum_{a'_1a'_2} \phi^{a(f)}_{a'_1}(S^a(n), \vec{r}_{2b})\phi^{f(b)}_{a'_2}(S^f(n), \vec{r}_{1b}), \quad \text{(7.C.11)}$$

where $(a_1, a_2) \equiv F$ and $(a'_1, a'_2) \equiv f$ span, as mentioned above, shells in nuclei B and a respectively.

Inserting (A9) and (A10) in (A4) one can show that

$$a^{(1)}(\infty) = a^{(NO)}(\infty). \quad \text{(7.C.12)}$$

It can be demonstrated that within the present approximation, $Im\ \tilde{a}^{(2)} = 0$, and that

$$\tilde{a}^{(2)}(\infty) = \frac{1}{i\hbar}\int_{-\infty}^{\infty} dt(\psi^b\psi^B, (V_{bB} - < V_{bB} >)e^{i\sigma_1}\psi^f\psi^F > \times \exp[\frac{i}{\hbar}(E^{bB} - E^{fF})t + \gamma_1(t)]$$
$$\times\frac{1}{i\hbar}\int_{-\infty}^{\infty} dt'(\psi^f\psi^F, (V_{fF} - < V_{fF} >)e^{i\sigma_2}\psi^a\psi^A) \times \exp[\frac{i}{\hbar}(E^{fF} - E^{aA})t' + \gamma_2(t')] \quad \text{(7.C.13)}$$

The total absolute differential cross section 7.C.2, where $P = |a(\infty)|^2 = |\tilde{a}^{(2)}|^2$, is then equal to the product of two one-particle transfer cross sections (see Fig. A3), associated with the (virtual) reaction channels

$$\alpha \equiv a + A \rightarrow f + F \equiv \gamma, \quad \text{(7.C.14)}$$

and

$$\gamma \equiv f + F \rightarrow b + B \equiv \beta. \quad \text{(7.C.15)}$$

In fact, Eq.(7.C.13) involves no time ordering and consequently the two processes above are completely independent of each other. This result was expected because being $v_{12} = 0$, the transfer of one nucleon cannot influence, aside form selecting the initial state for the second step, the behaviour of the other nucleon.

7.C.2 Strong correlation (cluster) limit

The second limit to be considered is the one in which the correlation betwen the two nucleons is so strong that (see Fig. A2)

$$S^a(2n) \approx S^a(n) >> S^f(n), \quad \text{(7.C.16)}$$

and

$$S^B(2n) \approx S^B(n) >> S^F(n). \quad \text{(7.C.17)}$$

That is, the magnitude of the one-nucleon separation energy is strongly modified by the pair breaking.

There is a different , although equivalent way to express (7.C.3) which is the more convenient to discuss the strong coupling limit. In fact, making use of the post-prior representation one can write

$$a^{(2)}(t) = \tilde{a}^{(2)}(t) - a^{(NO)}(t) = \frac{1}{i\hbar}\int_{-\infty}^{\infty} dt(\psi^b\psi^B, (V_{bB} - < V_{bB} >)e^{i\sigma_1}\psi^f\psi^F)$$
$$\times\exp[\frac{i}{\hbar}(E^{bB} - E^{fF})t + \gamma_1(t)] \times$$
$$\frac{1}{i\hbar}\int_{-\infty}^{t} dt'(\psi^f\psi^F, (V_{aA} - < V_{aA} >)\psi^a\psi^A) \times \exp[\frac{i}{\hbar}(E^{fF} - E^{aA})t' + \gamma_2(t')]. \quad \text{(7.C.18)}$$

The relations (7.C.16), (7.C.17) imply

$$E^{fF} - E^{aA} = S^a(n) - S^F(n) >> \frac{\hbar}{\tau}, \quad \text{(7.C.19)}$$

where τ is the collision time. Consequently the real part of $a^{(2)}(\infty)$ vanishes exponentially with the $Q-$value of the intermediate transition, while the imaginary part vanishes inversely proportional to this energy. One can thus write,

$$Re\ a^{(2)}(\infty) \approx 0, \quad \text{(7.C.20)}$$

and

$$a^{(2)}(\infty) \approx \frac{1}{i\hbar}\frac{\tau}{< E^{fF} > - E^{bB}} \sum_{fF} (\psi^b\psi^B, (V_{bB} - < V_{bB} >)\psi^f\psi^F)_{t=0} \times (\psi^f\psi^F, (V_{aA} - < V_{aA})\psi^a\psi^A)_{t=0}, \tag{7.C.21}$$

where one has utilized the fact that $E^{bB} \approx E^{aA}$. For $v_{12} \to \infty$, $(< E^{fF} > - E^{bB}) \to \infty$ and, consequently,

$$lim_{v_{12}\to\infty} a^{(2)}(\infty) = 0. \tag{7.C.22}$$

Thus the total two-nucleon transfer amplitude is equal, in the strong coupling limit, to the amplitude $a^{(1)}(\infty)$.

Summing up, only when successive transfer and non-orthogonal corrections are included in the description of the two-nucleon transfer process, does one obtain a consistent description of the process, which correctly converges to the weak and strong correlation limiting values.

Appendix 7.D Simple numerical estimates of successive and simultaneous transfer amplitudes

Let us denote

$$H = T + V, \tag{7.D.1}$$

the total hamiltonian describing the nuclear system, where V is the nuclear two–body interaction.

The fact that the nuclear quantality parameter has a value of $Q \approx 0.4$ testifies to the validity of independent particle motion in nuclei. This is tantamount to saying that there exist a single–particle potential U, such that

$$\langle\Psi_0|U|\Psi_0\rangle \ll \langle\Psi_0|(V - U)|\Psi_0\rangle, \tag{7.D.2}$$

where Ψ_0 is the exact ground state wavefunction, that is, $H\Psi_0 = E_0\Psi_0$. One can the write 7.D.1 as

$$H = T + V_{eff}, \tag{7.D.3}$$

where

$$V_{eff} = U + (V - U). \tag{7.D.4}$$

Let us now consider a reaction in which two nucleons are transferred between target and projectile, that is,

$$a(= b + 2) + A \to b + B(= A + 2). \tag{7.D.5}$$

The transfer cross section is proportional to the square of the amplitude

$$\sqrt{\sigma} \sim \langle bB|V_{eff}|aA\rangle = \langle bB|U|aA\rangle + \langle bB|(V - U)|aA\rangle. \tag{7.D.6}$$

Let us assume that the transferred nucleons are e.g. two neutrons moving in time reversal states lying close t othe Fermi energy (Cooper pair). In this case it is natural to assume that the operative component of $(V - U)$ is the pairing interaction

$$V_p = -GP^\dagger P, \tag{7.D.7}$$

where

$$P^{\dagger} = \sum_{\nu>0} a_{\nu}^{\dagger} a_{\tilde{\nu}}^{\dagger}, \tag{7.D.8}$$

is the pair operator, and

$$G \approx \frac{18}{A} \text{ MeV}, \tag{7.D.9}$$

is the pairing coupling constant for nucleons moving in an extended (2–3 major shell) configuration.

One can then write Eq. 7.D.6 as

$$\sqrt{\sigma} = \sqrt{\sigma_1} + \sqrt{\sigma_2}, \tag{7.D.10}$$

where

$$\sqrt{\sigma_1} \sim \langle Bb|U|aA\rangle \approx 2\left(\frac{|V_0|}{2}\right) O, \text{ (SUCC+NO)} \tag{7.D.11}$$

and

$$\begin{aligned} \sqrt{\sigma_2} \sim & \langle Bb|V-U|aA\rangle = \langle Bb|H_p|aA\rangle \\ & \approx GU(b)V(B) \approx \frac{G}{2}, \text{ (PAIRING)} \end{aligned} \tag{7.D.12}$$

In the process described by the transfer amplitude $\langle Bb|U|aA\rangle$, one nucleon is transferred under the effect of the single–particle potential of depth $V_0 (\approx -50$ MeV) while, simultaneously, the second nucleon moves over from a single–particle orbit centered around b to one centered around A profiting of the non–orthogonality of the corresponding wavefunctions $\varphi^{(b)}(r_{1b})$ and $\varphi^{(A)}(r_{1A})$. Within this context, it is then natural that O stands for the overlap between these two wavefunctions, that is, (see below simple estimate of O),

$$O = \langle \varphi^{(b)} | \varphi^{(A)} \rangle \approx 0.3 \times 10^{-2}, \tag{7.D.13}$$

and that (7.D.11) is known as the sum of the simultaneous plus non–orthogonality contributions to the two–nucleon transfer amplitude. Of notice that the prefactor 2 in (7.D.11) is associated with the fact that two nucleons can choose to jump from one system to the other through non–orthogonality while the factor $|V_0|/2$ is associated with the fact that transfer takes mainly place at the surface.

The term (7.D.12) corresponds to the simultaneous (t, p) transfer via the pairing two–body interaction V_p (see Eq. (7.D.7)), $U(A)$ and $V(B)$ being the product of two occupation amplitudes: $U(A)$ measures the availability of free single–particle orbitals around the Fermi energy in the target nucleus, while $V(B)$ reflects the degree of occupancy of levels in the residual system. Close to the Fermi energy $U(b)V(B) \approx (1/\sqrt{2})^2 = 1/2$, leading to the final expression of (7.D.12).

In keeping with the fact that the ratio of transfer amplitudes

$$\begin{aligned} \left(\frac{\sigma_1}{\sigma_2}\right)^{1/2} & \approx 2\frac{|V_0|}{2} \times O\frac{1}{G/2} \approx 2 \times A \times 10^{-2} \\ & \approx 2(A \approx 100), \end{aligned} \tag{7.D.14}$$

is larger that one, and that the correlation length of nuclear Cooper pairs ($\xi \approx \hbar v_F/2\Delta \approx$ 30 fm) is larger than nuclear dimensions, one can expect that the successive transfer of two nucleons under the influence of the single–particle field, can give an important contribution to the total transfer amplitude $\sqrt{\sigma}$. In other words, we expect the process

$$a(=b+2)+A \rightarrow f(=b+1)+F(A+1) \rightarrow b+B(=A+2) \tag{7.D.15}$$

gives a consistent contribution to $\sqrt{\sigma}$. The associated amplitude can be written as

$$\begin{aligned}\sqrt{\sigma_3} &\sim \sum_{fF} \frac{\langle bB|U|fF\rangle\langle fF|U|aA\rangle}{E_{aA}-E_{fF}} \\ &\approx \frac{(V_0/13)(V_0/13)}{\Delta E},\end{aligned} \tag{7.D.16}$$

the factor 1/7 appears in each of the steps (instead of 1/2, see (7.D.11)) in keeping with the fact that many other reactions channels and then, absorptive processes will take place at closer distance in two–step channels (of notice that 1/7 corresponds to $r = R_0 + 2.5a$).

Typical values of the energy denominator in (7.D.16) are $\Delta E = 30$ MeV for medium heavy nuclei lying along the stability valley.

Appendix 7.E Transfer amplitudes

Making use of a simplified expression for the elastic scattering amplitude, that is

$$\sqrt{\sigma_{el}} \sim \langle aA|U|aA\rangle, \tag{7.E.1}$$

one can calculate the transfer probabilities associated with the different processes discussed above, namely

$$P_i = \left(\frac{\sigma_i}{\sigma_{el}}\right) = \begin{cases} \left|\frac{\langle bB|U|aA\rangle}{\langle aA|U|aA\rangle}\right|^2 \approx O^2 \approx 0.9\times 10^{-5} \quad (i=1), \\ \left|\frac{\langle bB|V_n|aA\rangle}{\langle aA|U|aA\rangle}\right|^2 \approx \left(\frac{G}{2V_0}\right)^2 \approx 1.4\times 10^{-6} \quad (i=2), \\ \left|\frac{\langle bB|U|fF\rangle\langle fF|U|aA\rangle}{\Delta E\langle aA|U|aA\rangle}\right|^2 \approx \left(\frac{V_0}{170\Delta E}\right)^2 \approx 0.96\times 10^{-4} \quad (i=3). \end{cases} \tag{7.E.2}$$

Because all these probabilities are small, one can write

$$\sigma_i = P_i\sigma_{el}, \tag{7.E.3}$$

where

$$\begin{aligned}\sigma_{el} &= \left(\frac{\mu_\alpha}{2\pi\hbar^2}\right)^2 |\langle aA|U|aA\rangle|^2 \\ &\approx \left(\left(\frac{\mu_\alpha}{2\pi\hbar^2}\right)(V_0)\right)^2 U_0^2 \\ &\approx \left(1.8\text{MeV}^{-1}\text{fm}\right)^2 (50\text{MeV})^2 \\ &\approx (90\text{fm})^2 = 0.8\text{b}\end{aligned} \tag{7.E.4}$$

where use has been made of the effective volume associated with the reaction process (see)

$$V_{ol} \approx \frac{4\pi}{3} 3R^2 a \approx 12A^{2/3}\text{fm}^3 \approx 260\text{fm}^3, \tag{7.E.5}$$

as well as of $\frac{\mu_\alpha}{2\pi\hbar^2}$ factor of the typical two–nucleon transfer reaction ^{120}Sn+p$\rightarrow$^{118}Sn+t, that is (see),

$$\frac{\sqrt{\mu_\alpha\mu_\beta}}{2\pi\hbar^2} \approx \frac{\sqrt{3}M}{2\pi\hbar^2} \approx \frac{1}{145}\text{MeV}^{-1}\text{fm}^{-2}. \tag{7.E.6}$$

Summing up, one can write

$$\sigma_i = P_i 0.8\text{b}. \tag{7.E.7}$$

Making use of (7.E.2) one obtains

$$\sigma_i = \begin{cases} 0.7\times10^{-2}\text{mb} & (i=1), \\ 1.1\times10^{-3}\text{mb} & (i=2), \\ 8\,\text{mb} \quad (i=1), & (i=3). \end{cases} \tag{7.E.8}$$

These numbers, although worked out for A=100 can be rescaled in connection with the reaction ^{11}Li$(p,t)^9$Li(gs), in which case, microscopic calculation lead to $d\sigma_1(\theta = 60°)/d\Omega \approx 0.01$mb/sr and $d\sigma_3(\theta = 60°)/d\Omega \approx 5$mb/sr.

Appendix 7.F Inelastic scattering following two–particle transfer: final state interaction

This subject is qualitatively discussed in connection with the ^{11}Li$(p,t)^9$Li$(1/2^-; 2.69)$, but of course is a general question, also in connection with the validity of considering perturbation theory instead of coupled channels.

In keeping with the fact that the first excited state of ^{9}Li can be viewed as

$$|^9\text{Li}(1/2^-; 2.69\text{MeV})\rangle \approx |2^+\left(^8\text{Be}\otimes p_{3/2}(\pi)\right)_{1/2^-}\rangle, \tag{7.F.1}$$

this state can, in principle, be excited in a two–step process, namely

$$\text{gs}\ (^{11}\text{Li})+\text{t} \longrightarrow \ \text{gs}\ (^9\text{Li})+\text{p} \longrightarrow \ 1/2^-\ (^9\text{Li})+\text{p}. \tag{7.F.2}$$

Let us calculate the probability associated with the inelastic scattering of the lowest 2^+ of ^{8}Li. In this case, we are interested in the component of $V-U$ corresponding to $\delta U_C = -KF\alpha = -R_0\frac{\partial U}{\partial r}\beta_L$, namely the field associated with the inelastic excitation of multipole vibrations. Making use of the Saxon–Woos potential one obtains

$$R_0\frac{\partial U}{\partial r} = \frac{R_0}{a}\frac{\exp\left(\frac{(r-R_0)}{a}\right)}{\left(1+\exp\left(\frac{(r-R_0)}{a}\right)\right)^2}. \tag{7.F.3}$$

In keeping with the fact that

$$\left\langle R_0\left.\frac{\partial U}{\partial r}\right|_{r=R_0}\right\rangle \approx \left\langle\frac{R_0U_0}{a}\right\rangle \approx 1.2U_0\text{MeV} \approx -60\text{MeV} \tag{7.F.4}$$

and that the main contributions of surface dominated reactions processes is estimated to arize from distances of the order of $r \approx R_0 + 2.5a$, one obtains

$$\begin{aligned}\left\langle\frac{R_0}{a}\frac{e^{2.5}U_0}{(1+\exp 2.5)^2}\right\rangle &= \left\langle\frac{R_0U_0}{a}\right\rangle\frac{e^{2.5}}{(1+\exp 2.5)^2} \\ &\approx 1.2U_0\times0.7\times10^{-1} = 0.84\times10^{-1}U_0.\end{aligned} \tag{7.F.5}$$

Thus

$$\langle bB^*|\delta U_C|bB\rangle \approx 0.84\times 10^{-1}U_0\beta_L. \tag{7.F.6}$$

Consequently

$$\begin{aligned} P_{inel} &\approx \left|\frac{\langle bB^*|\delta U_C|aA\rangle}{\langle aA|U|aA\rangle}\right|^2 = \left(0.84\times 10^{-1}\beta_L\right)^2 \\ &\approx 0.7\times 10^{-2}\beta_L^2. \end{aligned} \tag{7.F.7}$$

In keeping with the fact that the β_L associated with the lowest 2^+ vibrational states of the Sn–isotopes and of ^{8}are ≈ 0.1 and ≈ 1 respectively one can write

$$P_{inel} = \begin{cases} 0.7\times 10^{-4} & \text{(Sn–isotopes)}, \\ 0.7\times 10^{-2} & (^{11}\text{Li}). \end{cases} \tag{7.F.8}$$

Making use of the results collected in (7.E.2),

$$\begin{aligned} \sqrt{P(p,t)} &= \sqrt{P_1} + \sqrt{P_2} + \sqrt{P_3} \\ &= \sqrt{0.9\times 10^{-1}} + \sqrt{1.4\times 10^{-6}} + \sqrt{0.96\times 10^{-4}} \\ &\approx 3\times 10^{-3} + 1.2\times 10^{-3} + 0.98\times 10^{-2} \\ &\approx 1.4\times 10^{-2}. \end{aligned} \tag{7.F.9}$$

Thus

$$P((p,t)\otimes P(\text{inel})) = P(p,t)P(\text{inel}) = \begin{cases} 2\times 10^{-4}\times 10^{-4}\approx 10^{-8} & \text{(Sn)}, \\ 2\times 10^{-4}\times 10^{-2}\approx 10^{-6} & (^{11}\text{Li}), \end{cases} \tag{7.F.10}$$

in overall agreement with th result of microscopic calculations (for ^{11}Li).

Appendix 7.G Simple estimate O

The nuclear density can be parametrized according to

$$\rho(r) = \frac{\rho_0}{1+\exp\left(\frac{r-R_0}{a}\right)}. \tag{7.G.1}$$

Let us calculate this function for

$$r = R_0 + 3a, \tag{7.G.2}$$

that is

$$\rho(r = R_0 + 3a) = \frac{\rho_0}{1+\exp 3} = 5\times 10^{-2}\rho_0. \tag{7.G.3}$$

In other words, we assume that the main transfer takes place from densities of the order of 5% saturation density

$$O \approx \frac{\rho_A(R_0^A + 3a)\rho_a(R_0^a + 3a)}{\rho_0^2} = 25\times 10^{-4} \approx 0.3\times 10^{-2}. \tag{7.G.4}$$

another estimate

$$r = R_0 + 2.5a, \tag{7.G.5}$$

for which

$$\rho(r = R_0 + 2.5a) = \frac{\rho_0}{1 + \exp 2.5} \approx 0.76 \times 10^{-1}\rho_0, \tag{7.G.6}$$

leading to

$$O \approx 0.5 \times 10^{-2}. \tag{7.G.7}$$

Appendix 7.H Simple estimate of $\frac{(\mu_\alpha\mu_\beta)^{1/2}}{2\pi\hbar^2}$.

Let us do it for the case of ^{120}Sn+p $\longrightarrow$ ^{118}Sn+t. In this case

$$\begin{aligned} \mu_\alpha &= \frac{120}{121}M \approx M, \\ \mu_\beta &= \frac{118 \times 3}{121} \approx 2.9M. \end{aligned} \tag{7.H.1}$$

Thus

$$\begin{aligned} \frac{\sqrt{\mu_\alpha\mu_\beta}}{2\pi\hbar^2} &\approx \frac{\sqrt{3}M}{2\pi\hbar^2} = \frac{\sqrt{3}}{2\pi 40\text{MeV fm}^2} \\ &\approx \frac{1}{145} \times \text{MeV}^{-1} \times \text{fm}^{-2} \end{aligned} \tag{7.H.2}$$

Appendix 7.I Simple estimate of V_{ol}

In keeping with the assumption that transfer processes are expected to take place at the nuclear surface, the effective volume associate with such processes can be estimated to be

$$\begin{aligned} V_{ol} &= \frac{4\pi}{3}(R^3 - (R + a)^3) \\ &\approx \frac{4\pi}{3}3aR^2 \approx \frac{4\pi}{3}(2\text{fm})R^2 \\ &\approx \frac{8\pi}{3}(1.2A^{1/3})^2\text{fm}^3 \\ &\approx 1.2A^{1/3}\text{fm}^3 \approx 260\text{fm}^3 (A \approx 100) \end{aligned} \tag{7.I.1}$$

Appendix 7.J Calculation of the (p,t) strength function

An important component of the interaction which binds the dineutron halo of ^{11}Li to the core ^{9}Li is associated with the exchange, between the two neutrons of dipole (pigmy 1^- resonance of ^{11}Li) and quadrupole (2^+ mode of ^{9}Li) vibrations. Consequently, it is expected that resonant effects can be observed in the (p,t) strength function in which this neutrons are picked-up from ^{11}Li. In keeping with the fact that successive transfer plays a central role in the two-particle pick up process, the corresponding transfer amplitude can be written as

$$\begin{aligned} &\left\langle \chi^{(-)} \sum_{fF} \frac{\langle bB|U|fF\rangle\langle fF|U|aA\rangle}{E_{aA} - E_{fF}} \chi^{(+)} \right\rangle \\ &\approx \left\langle \frac{e^{iqr}}{\hbar\omega_L} \right\rangle \sim \left\langle e^{i[qr - \ln\tau/\omega_L)} \right\rangle \\ &\sim \cos\left(q(E_{CM})r - \ln\tau/\omega_L\right). \end{aligned} \tag{7.J.1}$$

Of notice that (p,t)-strength function measurements can be viewed as a frequency dependent single Cooper pair transfer, and thus in some way connected to ω-dependent Josephson supercurrent measurement.

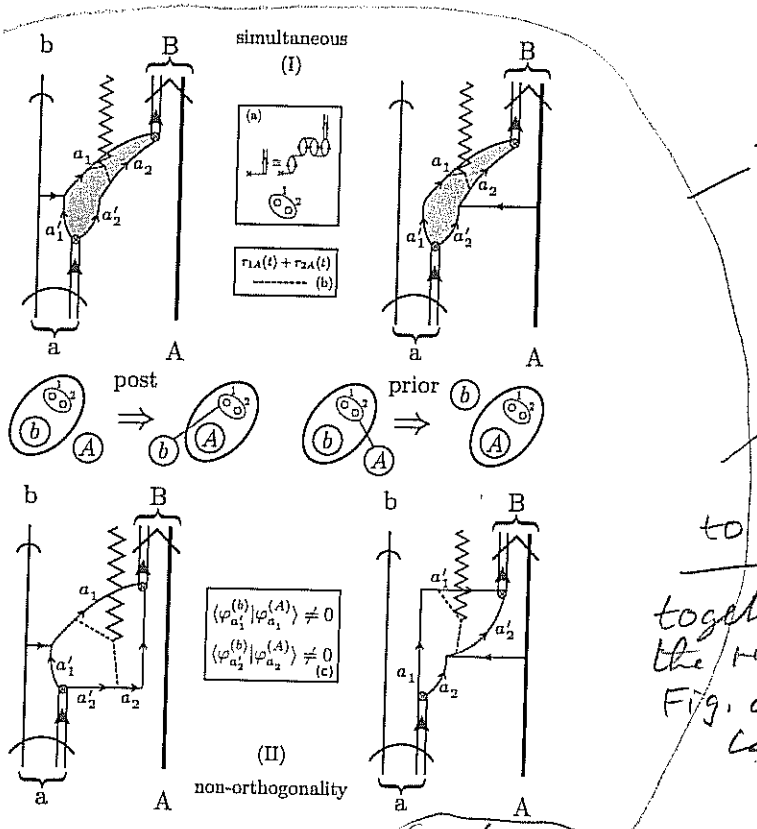


Fig 7.A.1

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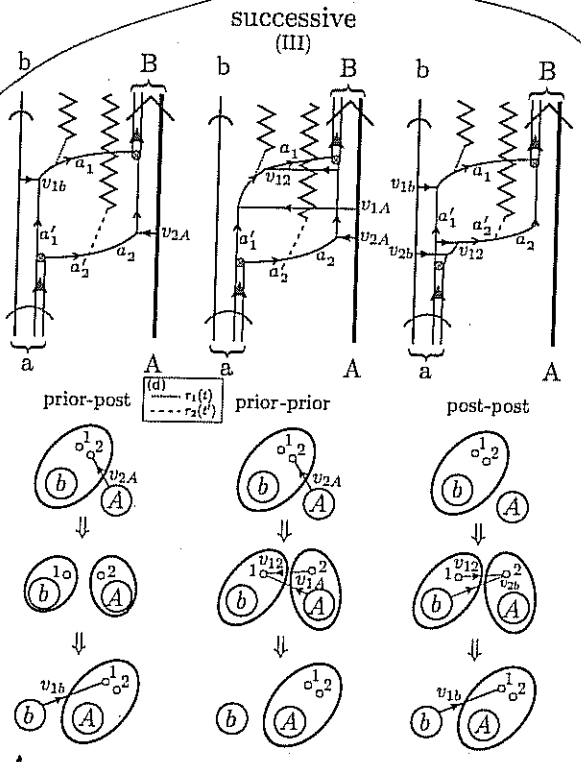

Figure 7.1: Graphical representation of the lowest order ((I),(II) and (III) first and second order in v respectively), two-nucleon transfer processes, which correctly converge to the strong-correlation (only simultaneous transfer), and to the independent-particle (only successive transfer) limits. The time arrow is assumed to point upwards: (I) Simultaneous transfer, in which one particle is transferred by the nucleon-nucleon interaction (note that $U(r) = \int d^3r' \rho(r') v(|\vec{r} - \vec{r}'|)$) acting either in the entrance $\alpha \equiv a + A$ channel (prior) or in the final $\beta \equiv b + B$ channel (post), while the other particle follows suit making use of the particle-particle correlation (grey area) which binds the Cooper pair (see upper inset labelled (a)), represented by a solid arrow on a double line, to the projectile (curved arrowed lines) or to the target (opened arrowed lines). The above argument provides the explanation why when e.g. v_{1b} acts on one nucleon, the other nucleon also reacts instantaneously. In fact a Cooper pair displays generalized rigidy (emergent property in gauge space). A crossed open circle represents the particle-pair vibration coupling. The associated strength, together with an energy denominator, determines the amplitude $X_{a'_1 a'_2}$ (cf. Table 1) with which the pair mode (Cooper pair) is in the (time reversered) two particle configuration $a'_1 a'_2$. In the transfer process, the orbital of relative motion changes, the readjustement of the corresponding trajectory mismatch being operated by a Galilean operator ($\exp[\vec{k} \cdot (\vec{r}_{1A}(t) + \vec{r}_{2A}(t))]$). This phenomenon, known as recoil process, is represented by a jagged line which provides simultaneous information on the two transferred nucleons (single time appearing as argument of both single-particle coordinates r_1 and r_2; see inset labeled (b)). In other words, information on the coupling of structure and reaction modes. (II) Non-orthogonality contribution. While one of the nucleons of the Cooper pairs is transferred under the action of v, the other goes, uncorrelatedly over, profiting of the non-orthogonality of the associated single-particle wavefunctions (see inset (c)). In other words of the non-vanishing values of the overlaps, as shown in the inset. (III) Successive transfer. In this case, there are two time dependences associated with the acting of the nucleon-nucleon interaction twice (see inset (d)).

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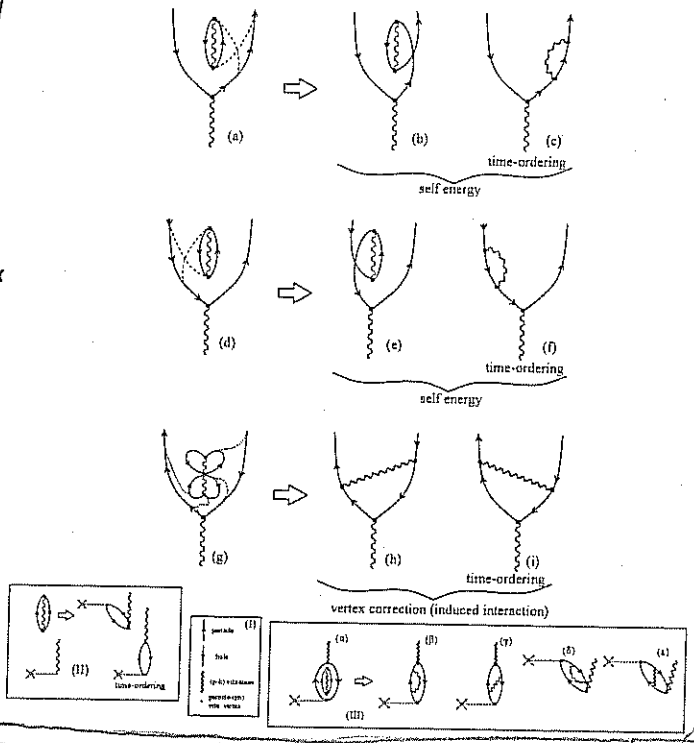

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correlated particle-hole excitation which in the shell model basis corresponds to a

formfactor and strength, i.e. transition density,

7.1.1

Figure 7.1.2: Nuclear field theory (NFT) diagrams corresponding to the lowest order medium polarization effects renormalizing the properties of a particle–hole collective mode (wavy line), linear combination of particle–hole ((up–going)–(down–going) arrowed lines) excitations calculated within the random phase approximation (RPA) in of a bare interaction, and leading to the particle-vibration coupling vertex (solid dot, see inset (I), bottom). The action of an external field on the zero point fluctuations (ZPF) of the vacuum (inset (II)), forces a virtual process to become real, leading to a collective vibration by annihilating a (virtual, spontaneous) particle–hole excitation (backwards RPA amplitude) or, in the time ordered process, by creating a particle–hole excitation which eventually, through the particle–vibration coupling vertex, correlates into the collective (coherent state; forwardsgoing amplitudes). Now, oyster–like diagrams associated with the vacuum ZPF can occur at any time (see inset (III)). Because the texture of the vacuum is permeated by symmetry rules (while one can violate energy conservation in a virtual state one cannot violate e.g. angular momentum conservation or the Pauli principle), the process shown in the inset III (α) leads, through Pauli principle correcting processes (exchange of fermionic arrowed lines) to self–energy (inset III (β), (δ)) and vertex corrections (induced p–h interaction; inset III (γ), (ε)) processes. The first ones are detailed in graphs (a)–(f), while the second ones in graphs (g)–(i). In keeping with the fact that the vibrational states can be viewed as a coherent state exhausting a large fraction of the EWSR (e.g. a Giant Resonance) for which the associated uncertainty relations in momentum and coordinate fulfills the absolute minimum consistent with quantum mechanics ($\Delta\alpha_{\lambda\mu}\Delta\pi_{\lambda\mu} = \hbar/2, \alpha_{\lambda\mu} = (\hbar\omega_\mu/2C_\lambda^{1/2})(\Gamma^\dagger_{\lambda\mu} + \Gamma_{\lambda\mu})$ being the (harmonic) collective coordinate, $\pi_{\lambda\mu}$ being the conjugate momentum; cf. e.g. [illegible]), there is a strong cancellation between the contribution of self–energy and vertex correction diagrams (P. F. Bortignon et al., [illegible]), implying small anharmonicities and long lifetimes ($\Gamma/E \ll 1$, where Γ is the width and E the centroid of the mode $|\lambda\mu\rangle = \Gamma^\dagger_{\lambda\mu}|0\rangle$, $(\hbar\omega_\lambda/2C_\lambda)^{1/2}$ being the ZPF amplitude (cf. e.g. Brink and Broglia (2005))).

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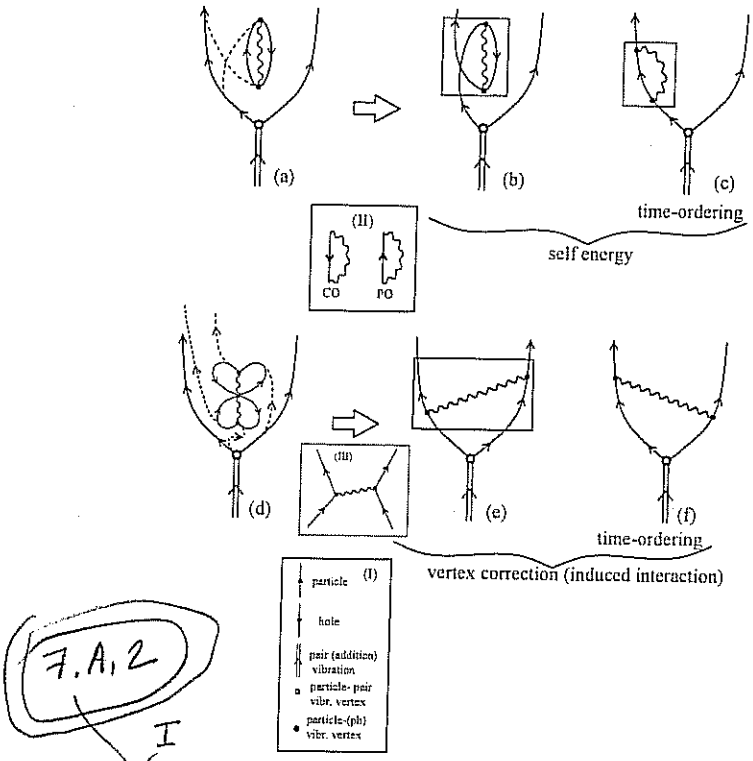

Figure 7.A.2: Pauli effects associated (p–h) ZPF dressing of pairing vibrational modes (see inset bottom-left) in terms of self-energy (graphs (a)–(c); correlation (CO) and polarization (PO) diagrams, inset II) and vertex correction (graphs (d)–(f); induced particle–particle interaction (inset (III)) processes.