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Let us examine the term

$$\sum_M (-1)^M \langle K\ \Lambda\ -M\ \mu | P\ \mu - M\rangle \left[Y^{l_f}(\hat{\mathbf{r}}_{A1}) Y^{l_{i1}}(\hat{\mathbf{r}}_{b1})\right]^K_M \left[Y^{l_f}(\hat{\mathbf{r}}'_{A2}) Y^{l_{i2}}(\hat{\mathbf{r}}'_{b2})\right]^P_{\mu-M}. \tag{7.2.167}$$

Making use of the relation

$$\langle l_1\ l_2\ m_1\ m_2 | L\ M_L\rangle = (-1)^{l_2-m_2}\sqrt{\frac{2L+1}{2l_1+1}}\langle L\ l_2\ -M_L\ m_2 | l_1\ -m_1\rangle, \tag{7.2.168}$$

the expression (7.2.168) is equivalent to,

$$(-1)^K\sqrt{\frac{2P+1}{2\Lambda+1}}\left\{\left[Y^{l_f}(\hat{\mathbf{r}}'_{A2}) Y^{l_{i2}}(\hat{\mathbf{r}}'_{b2})\right]^P \left[Y^{l_f}(\hat{\mathbf{r}}_{A1}) Y^{l_{i1}}(\hat{\mathbf{r}}_{b1})\right]^K\right\}^\Lambda_\mu. \tag{7.2.169}$$

We now recouple the term

$$\left[Y^{l_a}(\hat{\mathbf{r}}'_{aA}) Y^{l_a}(\hat{\mathbf{k}}_{aA})\right]^0_0 \left[Y^{l_b}(\hat{\mathbf{r}}_{bB}) Y^{l_b}(\hat{\mathbf{k}}_{bB})\right]^0_0, \tag{7.2.170}$$

arising from the partial wave expansion of the incoming and outgoing distorted waves to have,

$$((l_a l_a)_0 (l_b l_b)_0 | (l_a l_b)_\Lambda (l_a l_b)_\Lambda)_0 \left\{\left[Y^{l_a}(\hat{\mathbf{r}}'_{aA}) Y^{l_b}(\hat{\mathbf{r}}_{bB})\right]^\Lambda \left[Y^{l_a}(\hat{\mathbf{k}}_{aA}) Y^{l_b}(\hat{\mathbf{k}}_{bB})\right]^\Lambda\right\}^0_0. \tag{7.2.171}$$

The only term which does not vanish upon integration is

$$\frac{(-1)^{\Lambda-\mu}}{\sqrt{(2l_a+1)(2l_b+1)}}\left[Y^{l_a}(\hat{\mathbf{r}}'_{aA}) Y^{l_b}(\hat{\mathbf{r}}_{bB})\right]^\Lambda_{-\mu} \left[Y^{l_a}(\hat{\mathbf{k}}_{aA}) Y^{l_b}(\hat{\mathbf{k}}_{bB})\right]^\Lambda_\mu. \tag{7.2.172}$$

Again, the only term surviving

$$\left\{\left[Y^{l_f}(\hat{\mathbf{r}}'_{A2}) Y^{l_{i2}}(\hat{\mathbf{r}}'_{b2})\right]^P \left[Y^{l_f}(\hat{\mathbf{r}}_{A1}) Y^{l_{i1}}(\hat{\mathbf{r}}_{b1})\right]^K\right\}^\Lambda_\mu \left[Y^{l_a}(\hat{\mathbf{r}}'_{aA}) Y^{l_b}(\hat{\mathbf{r}}_{bB})\right]^\Lambda_{-\mu} \tag{7.2.173}$$

is

$$\frac{(-1)^{\Lambda+\mu}}{\sqrt{2\Lambda+1}}\left[\left\{\left[Y^{l_f}(\hat{\mathbf{r}}'_{A2}) Y^{l_{i2}}(\hat{\mathbf{r}}'_{b2})\right]^P \left[Y^{l_f}(\hat{\mathbf{r}}_{A1}) Y^{l_{i1}}(\hat{\mathbf{r}}_{b1})\right]^K\right\}^\Lambda \left[Y^{l_a}(\hat{\mathbf{r}}'_{aA}) Y^{l_b}(\hat{\mathbf{r}}_{bB})\right]^\Lambda\right]^0_0. \tag{7.2.174}$$

We now couple this last term with the term $\left[Y^{l_c}(\hat{\mathbf{r}}'_{cC})Y^{l_c}(\hat{\mathbf{r}}_{cC})\right]^0_0$, arising from the partial wave expansion of the Green function. That is,

$$
\begin{aligned}
&\left[\left\{\left[Y^{l_f}(\hat{\mathbf{r}}'_{A2})Y^{l_{i2}}(\hat{\mathbf{r}}'_{b2})\right]^P\left[Y^{l_f}(\hat{\mathbf{r}}_{A1})Y^{l_{i1}}(\hat{\mathbf{r}}_{b1})\right]^K\right\}^\Lambda\left[Y^{l_a}(\hat{\mathbf{r}}'_{aA})Y^{l_b}(\hat{\mathbf{r}}_{bB})\right]^\Lambda\right]^0_0\left[Y^{l_c}(\hat{\mathbf{r}}'_{cC})Y^{l_c}(\hat{\mathbf{r}}_{cC})\right]^0\\
&\quad=((l_al_b)_\Lambda(l_cl_c)_0|(l_al_c)_P(l_bl_c)_K)_\Lambda\left[\left\{\left[Y^{l_f}(\hat{\mathbf{r}}'_{A2})Y^{l_{i2}}(\hat{\mathbf{r}}'_{b2})\right]^P\left[Y^{l_f}(\hat{\mathbf{r}}_{A1})Y^{l_{i1}}(\hat{\mathbf{r}}_{b1})\right]^K\right\}^\Lambda\right.\\
&\quad\left.\left\{\left[Y^{l_a}(\hat{\mathbf{r}}'_{aA})Y^{l_c}(\hat{\mathbf{r}}'_{cC})\right]^P\left[Y^{l_b}(\hat{\mathbf{r}}_{bB})Y^{l_c}(\hat{\mathbf{r}}_{cC})\right]^K\right\}^\Lambda\right]^0_0=((l_al_b)_\Lambda(l_cl_c)_0|(l_al_c)_P(l_bl_c)_K)_\Lambda\\
&\quad\times((PK)_\Lambda(PK)_\Lambda|(PP)_0(KK)_0)_0\left\{\left[Y^{l_f}(\hat{\mathbf{r}}'_{A2})Y^{l_{i2}}(\hat{\mathbf{r}}'_{b2})\right]^P\left[Y^{l_a}(\hat{\mathbf{r}}'_{aA})Y^{l_c}(\hat{\mathbf{r}}'_{cC})\right]^P\right\}^0_0\\
&\quad\times\left\{\left[Y^{l_f}(\hat{\mathbf{r}}_{A1})Y^{l_{i1}}(\hat{\mathbf{r}}_{b1})\right]^K\left[Y^{l_b}(\hat{\mathbf{r}}_{bB})Y^{l_c}(\hat{\mathbf{r}}_{cC})\right]^K\right\}^0_0=((l_al_b)_\Lambda(l_cl_c)_0|(l_al_c)_P(l_bl_c)_K)_\Lambda\\
&\quad\times\sqrt{\frac{2\Lambda+1}{(2K+1)(2P+1)}}\left\{\left[Y^{l_f}(\hat{\mathbf{r}}'_{A2})Y^{l_{i2}}(\hat{\mathbf{r}}'_{b2})\right]^P\left[Y^{l_a}(\hat{\mathbf{r}}'_{aA})Y^{l_c}(\hat{\mathbf{r}}'_{cC})\right]^P\right\}^0_0\\
&\quad\times\left\{\left[Y^{l_f}(\hat{\mathbf{r}}_{A1})Y^{l_{i1}}(\hat{\mathbf{r}}_{b1})\right]^K\left[Y^{l_b}(\hat{\mathbf{r}}_{bB})Y^{l_c}(\hat{\mathbf{r}}_{cC})\right]^K\right\}^0_0.
\end{aligned}
\tag{7.2.175}
$$

Collecting all the contributions (including the constants and phases arising from the partial wave expansion of the distorted waves and the Green function), we get

$$
\begin{aligned}
T^{succ}_\mu=&(-1)^{j_f+j_{i1}}\frac{2048\pi^5\mu_{Cc}}{\hbar^2k_{Aa}k_{Bb}k_{Cc}}\sqrt{\frac{(2j_{i1}+1)}{(2\Lambda+1)(2j_f+1)}}\sum_{K,P}((l_f\tfrac{1}{2})_{j_f}(l_{i2}\tfrac{1}{2})_{j_{i2}}|(l_fl_{i2})_P(\tfrac{1}{2}\tfrac{1}{2})_0)_P\\
&\times((l_f\tfrac{1}{2})_{j_f}(l_{i1}\tfrac{1}{2})_{j_{i1}}|(l_fl_{i1})_K(\tfrac{1}{2}\tfrac{1}{2})_0)_K\;((j_{i1}j_f)_K(j_{i1}j_{i2})_\Lambda|(j_{i1}j_{i1})_0(j_fj_{i2})_P)_P\\
&\times\frac{(-1)^K}{(2K+1)\sqrt{2P+1}}\sum_{l_c,l_a,l_b}((l_al_b)_\Lambda(l_cl_c)_0|(l_al_c)_P(l_bl_c)_K)_\Lambda e^{i(\sigma_i^{l_a}+\sigma_f^{l_b})}i^{l_a-l_b}\\
&\times(2l_c+1)^{3/2}\left[Y^{l_a}(\hat{\mathbf{k}}_{aA})Y^{l_b}(\hat{\mathbf{k}}_{bB})\right]^\Lambda_\mu S_{K,P,l_a,l_b,l_c},
\end{aligned}
\tag{7.2.176}
$$

with (note that we have reduced the dimensionality of the integrals in the same fashion as for the L =0–angular momentum transfer calculation, see (7.2.132))

$$
\begin{aligned}
S_{K,P,l_a,l_b,l_c}=&\int r^2_{Cc}\,dr_{Cc}\,r^2_{b1}\,dr_{b1}\;\sin\theta\,d\theta\,v(r_{b1})u_{l_f}(r_{C1})u_{l_i}(r_{b1})\\
&\times\frac{s_{P,l_a,l_c}(r_{Cc})}{r_{Cc}}\frac{F_{l_b}(r_{Bb})}{r_{Bb}}\\
&\times\sum_M\langle l_c\;0\;l_b\;M|K\;M\rangle\left[Y^{l_f}(\hat{r}_{C1})Y^{l_{i1}}(\theta+\pi,0)\right]^K_M Y^{l_b}_{-M}(\hat{r}_{Bb}),
\end{aligned}
\tag{7.2.177}
$$

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and

$$s_{P,l_a,l_c}(r_{Cc}) = \int r'^2_{Cc}\, dr'_{Cc}\, r'^2_{A2}\, dr'_{A2}\, \sin\theta'\, d\theta'\, v(r'_{c2}) u_{l_f}(r'_{A2}) u_{l_i}(r'_{c2})$$
$$\times \frac{F_{l_a}(r'_{Aa})}{r'_{Aa}} \frac{f_{l_c}(k_{Cc}, r_<) P_{l_c}(k_{Cc}, r_>)}{r'_{Cc}} \tag{7.2.178}$$
$$\times \sum_M \langle l_c\ 0\ l_a\ M | P\ M \rangle \left[Y^{l_f}(\hat{r}'_{A2}) Y^{l_i}(\hat{r}'_{c2}) \right]^P_M Y^{l_a}_{-M}(\hat{r}'_{Aa}).$$

We have evaluated the transition matrix element for a particular projection μ of the initial angular momentum of the two transferred nucleons. If they are coupled to a core of angular momentum J_f to total angular momentum J_i, M_i, the fraction of the initial wavefunction with projection μ is $\langle \Lambda\ \mu\ J_f\ M_i - \mu | J_i\ M_i \rangle$, and the cross section will be

$$\frac{d\sigma}{d\Omega}(\hat{k}_{bB}) = \frac{k_{bB}}{k_{aA}} \frac{\mu_{aA}\mu_{bB}}{(2\pi\hbar^2)^2} \left| \sum_\mu \langle \Lambda\ \mu\ J_f\ M_i - \mu | J_i\ M_i \rangle T_\mu \right|^2. \tag{7.2.179}$$

For a non polarized incident beam,

(245) App A

$$\frac{d\sigma}{d\Omega}(\hat{k}_{bB}) = \frac{k_{bB}}{k_{aA}} \frac{\mu_{aA}\mu_{bB}}{(2\pi\hbar^2)^2} \frac{1}{2J_i + 1} \sum_{M_i} \left| \sum_\mu \langle \Lambda\ \mu\ J_f\ M_i - \mu | J_i\ M_i \rangle T_\mu \right|^2. \tag{7.2.180}$$

This would be the differential cross section for a transition to a definite final state M_f. If we do not measure M_f we have to sum for all M_f,

gregory

$$\frac{d\sigma}{d\Omega}(\hat{k}_{bB}) = \frac{k_{bB}}{k_{aA}} \frac{\mu_{aA}\mu_{bB}}{(2\pi\hbar^2)^2} \frac{1}{2J_i + 1} \sum_\mu |T_\mu|^2 \sum_{M_i, M_f} \left| \langle \Lambda\ \mu\ J_f\ M_f | J_i\ M_i \rangle \right|^2. \tag{7.2.181}$$

The sum over M_i, M_f of the Clebsh–Gordan coefficients gives $(2J_i + 1)/(2\Lambda + 1)$ (see ??) One then gets,

$$\frac{d\sigma}{d\Omega}(\hat{k}_{bB}) = \frac{k_{bB}}{k_{aA}} \frac{\mu_{aA}\mu_{bB}}{(2\pi\hbar^2)^2} \frac{1}{(2\Lambda + 1)} \sum_\mu |T_\mu|^2, \tag{7.2.182}$$

where one can write

$$T_\mu = \sum_{l_a, l_b} C_{l_a, l_b} \left[Y^{l_a}(\hat{k}_{aA}) Y^{l_b}(\hat{k}_{bB}) \right]^\Lambda_\mu$$
$$= \sum_{l_a, l_b} C_{l_a, l_b} i^{l_a} \sqrt{\frac{2l_a + 1}{4\pi}} \langle l_a\ l_b\ 0\ \mu | \Lambda\ \mu \rangle Y^{l_b}_\mu(\hat{k}_{bB}). \tag{7.2.183}$$

7,2

Note that (7.2.182) takes into account only the spins of the heavy nucleus. In a (t, p) or (p, t) reaction, we have to sum over the spins of the proton and of the triton and divide by 2. If a spin orbit term is present in the optical potential, the sum yields the combination of terms shown in section (??)

$$\frac{d\sigma}{d\Omega}(\hat{k}_{bB}) = \frac{k_{bB}}{k_{aA}} \frac{\mu_{aA}\mu_{bB}}{(2\pi\hbar^2)^2} \frac{1}{2(2\Lambda + 1)} \sum_\mu |A_\mu|^2 + |B_\mu|^2. \tag{7.2.184}$$

note different factors in Eq. (7,2,40)

parte de las correcciones están en la versión ~~appendix~~ apendice

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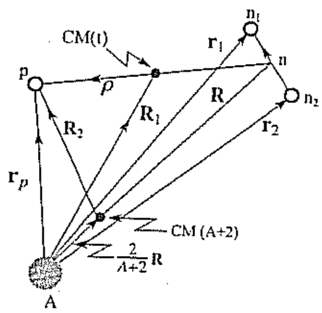
Appendix 7.A ZPF and Pauli principle at the basis of medium polarization effects: self–energy, vertex corrections and induced interaction

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.J.2 and 7.J.3); within the present context, see also Schrieffer (1964).

parcialmente corregido

Appendix 7.B Coherence and effective formfactors



$\mathbf{r} = \mathbf{r}_1 - \mathbf{r}_2$ (relative distance between the neutrons)

$\mathbf{R} = \frac{\mathbf{r}_1+\mathbf{r}_2}{2}$ (coord. of the CM of the dineutron)

$\rho = \mathbf{r}_p - \frac{\mathbf{r}_1+\mathbf{r}_2}{2} = \mathbf{r}_p - \mathbf{R}$ (distance between the CM of the dineutron and the proton)

$\mathbf{R}_2 = \mathbf{r}_p - \frac{\mathbf{r}_1+\mathbf{r}_2}{A+2}$ (distance of the proton from the CM of the system A+2)

$\mathbf{R}1 = \frac{\mathbf{r}_p+\mathbf{r}_1+\mathbf{r}_2}{3}$ (coord. of the CM of the triton)

Figure 7.B.1 out

coordinate system used in the calculation of the transfer amplitude

In what follows we shall work ~~on~~ a simplified derivation of the simultaneous two–nucleon transfer amplitude, within the framework of first order DWBA specially suited to discuss correlation aspects of pair transfer in general, and of the associated effective formfactors in particular.

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

a

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

ℓ

i.e $(sM'_s\, s'_f M'_{s_f}|s_iM_{s_i})\chi^{s'_f}_{M'_{s_f}}(\sigma_\beta)$

while

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

Making use of the above equation one can define the spectroscopic amplitude B as

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

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

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

is the spin function of the two–neutron system.

A convenient description of the intrinsic degrees of freedom of the triton is obtained by using a wavefunction symmetric in the coordinates of all particles, i.e.

$$\begin{aligned}\phi_t^{L=0}(\sum_{i<j}|\vec{r}_i-\vec{r}_j|) &= N_t\, e^{[(r_1-r_2)^2+(r_1-r_p)^2+(r_2-r_p)^2]}\\ &= \phi_{000}(\vec{r})\phi_{000}(\vec{\rho}),\end{aligned} \tag{7.B.6}$$

where

$$\phi_{000}(\vec{r})=R_{nl}(\nu^{1/2}r)Y_{lm}(\hat{r}).$$

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 (cf. Fig. 7.B.1)

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

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

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

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

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

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

$$T = \int d\xi_A\, d\vec{r}_1\, d\vec{r}_2\, d\vec{r}_p \chi_p^{(-)}(\vec{R}_2)\psi_\beta^*(\xi_{A+2}, \sigma_\beta) V_\beta \psi_\alpha(\xi_A, \sigma_\alpha, \sigma_\beta)\psi_t^{(+)}(\vec{R}_1), \tag{7.B.8}$$

where the final state effective interaction $V_\beta(\rho)$ is assumed to depend only on the distance ρ between the center of mass of the di–neutron and of the proton. 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 such If we ~~assume~~ that both $\phi_{j_1}(\vec{r}_1)$ and $\phi_{j_2}(\vec{r}_2)$ are harmonic oscillator wave functions (used as a basis to expand the Saxon-Woods single-particle wavefunctions) ~~this~~ transformation can ~~easily~~ be carried with the aid of tha Moshinsky brackets. If $|n_1l_1, n_2l_2; \lambda\mu\rangle$ is a complete 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

$$\begin{aligned} |n_1l_1, n_2l_2; \lambda\mu\rangle &= (\sum_{nlNL} |nl, NL; \lambda\mu\rangle\langle nl, NL; \lambda\mu|)|n_1l_1, n_2l_2; \lambda\mu\rangle \\ &= \sum_{nlNL} |nl, NL; \lambda\mu\rangle\langle nl, NL; \lambda\mu|n_1l_1, n_2l_2; \lambda\rangle . \end{aligned} \tag{7.B.9}$$

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

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

The coefficients $\langle nl, NL, L|n_1l_1, n_2l_2, L\rangle$ are tabulated and were first discussed by ~~M.~~ ~~Moshinsky in Nucl. Physics, 13 (1959) 104.~~ Moshinsky, 1959

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

$$\begin{aligned}
\psi^{J_f}_{M_f}(\xi_{A+2}) &= \sum_{\substack{n_1l_1j_1\\n_2l_2j_2\\JJ'_i}} B(n_1l_1j_1, n_2l_2j_2; JJ'_iJ_f)\left[\phi^J(j_1j_2)\phi^{J'_i}(\xi_A)\right]^{J_f}_{M_f}\\
&= \sum_{\substack{n_1l_1j_1\\n_2l_2j_2}} \sum_{JJ'_i} B(n_1l_1j_1, n_2l_2j_2; JJ'_iJ_f)\\
&\times \sum_{M_JM'_{J_i}} \langle JM_JJ'_iM_{J_i}|J_fM_{J_f}\rangle \psi^{J'_i}_{M'_{J_i}}(\xi_A)\\
&\times \sum_{LS'} \langle S'LJ|j_1j_2J\rangle \sum_{M_LM'_S} \langle LM_LS'M'_S|JM_J\rangle \chi^{S'}_{M'_S}(\sigma_\alpha)\\
&\times \sum_{nlN\Lambda} \langle nl, N\Lambda, L|n_1l_1, n_2l_2, L\rangle\\
&\times \sum_{m_lM_\Lambda} \langle lm_l\Lambda M_\Lambda|LM_L\rangle \phi_{nlm_l}(\vec r)\phi_{N\Lambda M_\Lambda}(\vec R) .
\end{aligned} \tag{7.B.11}$$

(non-orthogonality effect)

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

$V(\rho)$

$$\begin{aligned}
(\psi^{J_i}_{M_{J_i}}, V'_\beta(\rho)\psi^{J'_i}_{M'_{J_i}}) &= \delta(J_i, J'_i)\delta(M_{J_i}, M_{J'_i})V_\beta(\rho) \; ,\\
(\chi^S_{M_S}(\sigma_\alpha), \chi^{S'}_{M_{S'}}(\sigma_\alpha)) &= \delta(S,S')\delta(M_S, M_{S'}) \; ,\\
(\chi^{S_f}_{M_{S_f}}(\sigma_\beta), \chi^{S'_f}_{M_{S'_f}}(\sigma_\beta)) &= \delta(S_f, S'_f)\delta(M_{S_f}, M_{S'_f}) \; .
\end{aligned} \tag{7.B.14}$$

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

$V(\rho)$

$$\begin{aligned}
T = &\sum_{\substack{n_1l_1j_1\\n_2l_2j_2}} \sum_{JM_J} \sum_{nN} \sum_{S} B(n_1l_1j_1, n_2l_2j_2; JJ'_iJ_f)\\
&\times \langle JM_JJ_iM_{J_i}|J_fM_{J_f}\rangle\langle SLJ|j_1j_2J\rangle\\
&\times \langle LM_LSM_S|JM_J\rangle\langle n0, NL, L|n_1l_1, n_2l_2, L\rangle\\
&\times \langle SM_SS_fM_{S_f}|S_iM_{S_i}\rangle\Omega_n\\
&\times \int d\vec R\, d\vec r_p\, \chi^{(+)*}_t(\vec R_1)\phi^*_{NLM_L}(\vec R)V_\beta(\rho)\phi_{000}(\vec\rho)\chi^{(+)}_i(\vec R_1) \; ,
\end{aligned} \tag{7.B.15}$$

where we have approximated V_β by an effective interaction depending only on $\rho = |\vec\rho|$.

We now define the effective two–nucleon transfer form factor as

$$u^{J_iJ_f}_{LSJ}(R) = \sum_{n_1l_1j_1} B(n_1l_1j_1, n_2l_2j_2; JJ_iJ_f)\langle SLJ|j_1j_2J\rangle$$
$$\langle n0, NL, L|n_1l_1, n_2l_2; L\rangle \Omega_n R_{nL}(R) \tag{7.B.16}$$

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

$$T = \sum_J \sum_L \sum_S (JM_JJ_iM_{J_i}|J_fM_{J_f})(SM_SS_fM_{S_f}|S_iM_{S_i})(LM_LSM_S|JM_J)$$
$$\times \int d\vec{R}\, d\vec{r}_p \chi^{*(-)}_p(\vec{R}_2) u^{J_iJ_f}_{LSJ}(R) Y^*_{LM_L} V(\rho)\phi_{000}(\vec{\rho})\chi^{(+)}_t(\vec{R}_1) \tag{7.B.17}$$

Because the di–neutron has $S = 0$, we have that 6.F, Eq. (6.F.15)

$$\cancel{(LM_L00|JM_J)} = \delta(J,L)\delta(M_L,M_J) \tag{7.B.18}$$

and the summations over S and L disappear from eq. (7.B.17). Let us now make also here, as done in App. for one particle transfer reactions, the zero range approximation, that is, one-particle

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

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. Within this approximation (cf. Fig. 6.F.1)

~~From eqs. (??) we obtain~~

$$\vec{R} = \vec{R}_1 = \vec{r},$$
$$\vec{R}_2 = \frac{A}{A+2}\vec{R}, \tag{7.B.20}$$

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

lower case

$$T = D_0 \sum_L (LM_LJ_iM_{J_i}|J_fM_{J_f})$$
$$\times \int d\vec{R}\chi^{*(-)}_p\left(\frac{A}{A+2}\vec{R}\right) u^{J_iJ_f}_L(R) Y^*_{LM_L}(\hat{R})\chi^{(+)}_t(\vec{R}) \tag{7.B.21}$$

Examples of two-nucleon transfer form factors are given in Figs. 7.B.2, 7.B.3 and 7.B.4 Broglia and Riedel 1967

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~~ in zero-range approximation are,

lower case

$$\Delta S = 0$$
$$\Delta J = \Delta L = L \tag{7.B.22}$$
$$\Delta\pi = (-1)^L$$

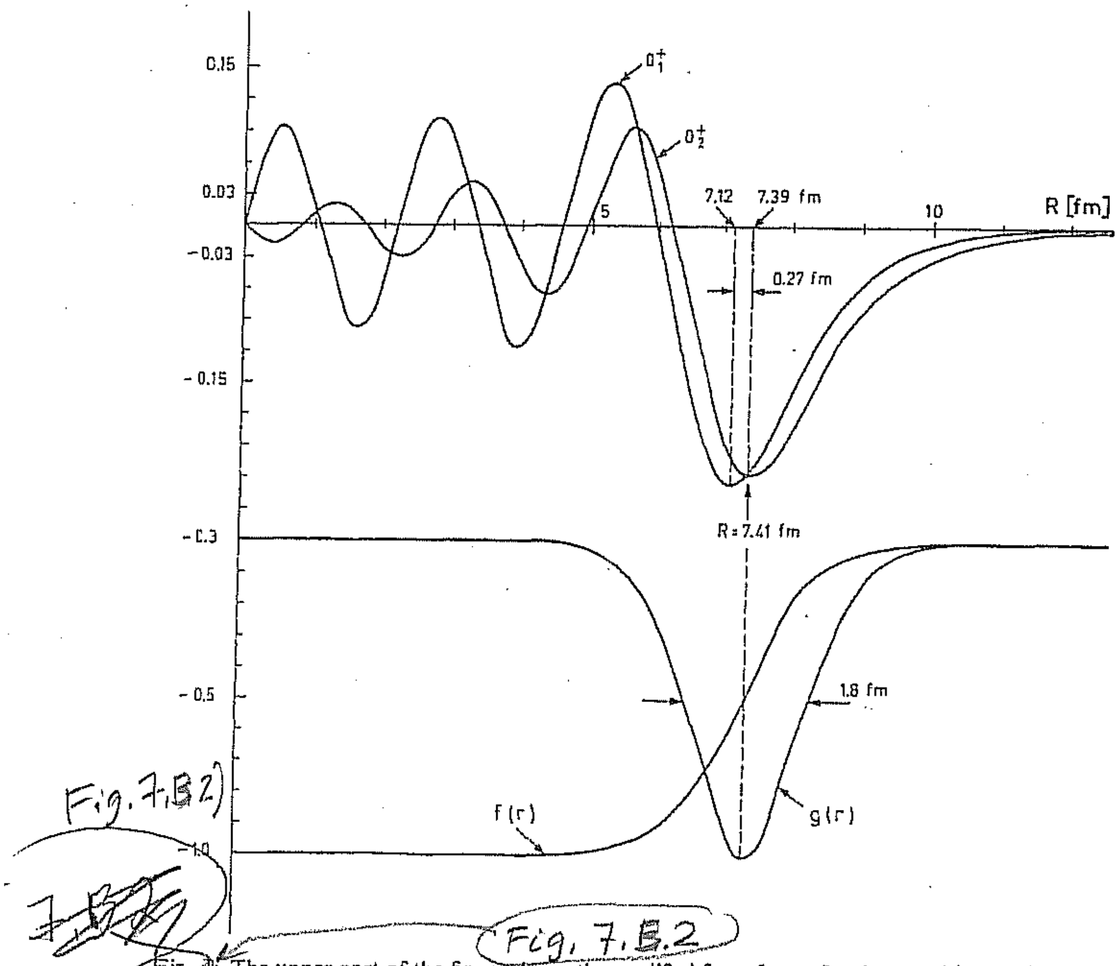
among other things,

i.e. only normal parity states are excited.

The integral appearing in eq. (7.B.21) has the same structure as the DWBA integral appearing in Eq. (6.F.16) ~~(??)~~ derived for 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 (6.F.1) ~~(??)~~ is a single–particle bound state wave function, $u^{J_iJ_f}_L$ is a coherent summation over the center of mass states of motion of the two transferred neutrons. In other words, an effective quantity (function). It is of notice ~~that~~ this difference essentially vanishes, when one considers dressed particles resulting from ~~the~~ coupling to collective motion, and leading to ω-dependent effective masses.

(i) a pure shell-model configuration $(p_{\frac{1}{2}})^{-2}$ for the ground state of ^{206}Pb, (ii) a configuration mixing caused by pairing a residual interaction and (iii) the same model as in (ii) including ground state correlations.
The results show clearly that different assumptions about the structure of one definite nuclear state lead to almost the same shape of the f_L functions inside the nucleus. The



Fig, 7.B.2

Fig. 8. The upper part of the figure shows the modified form factor for the transition to the ground state (0^{1+}) and the pairing vibrational state (0^{2+}) at 4.87 MeV. Both curves are matched with appropriate Hankel functions. In the lower part the form factors of the real ($f(r)$) and the imaginary ($g(r)$) part of the optical potential are given in the same scale for the radius.

form factor calculated by means of the shell-model wave functions of True and Ford[13]) is very similar to that of model (ii). The main difference of the three bound states is the magnitude of the maximum around the nuclear surface. Then, coherence effects only affect the degree of collectiveness of the form factor but not its shape.

Let us now discuss some details of the asymptotic behaviour of the form factors.

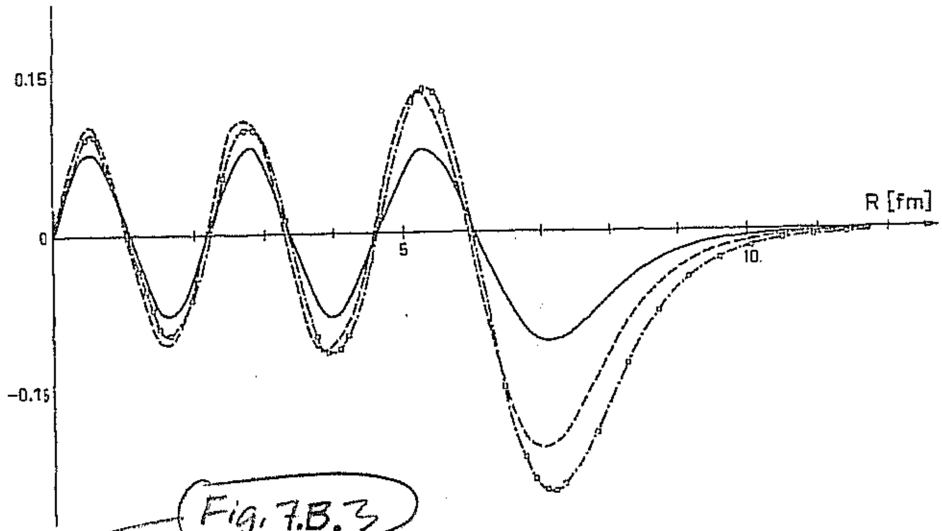
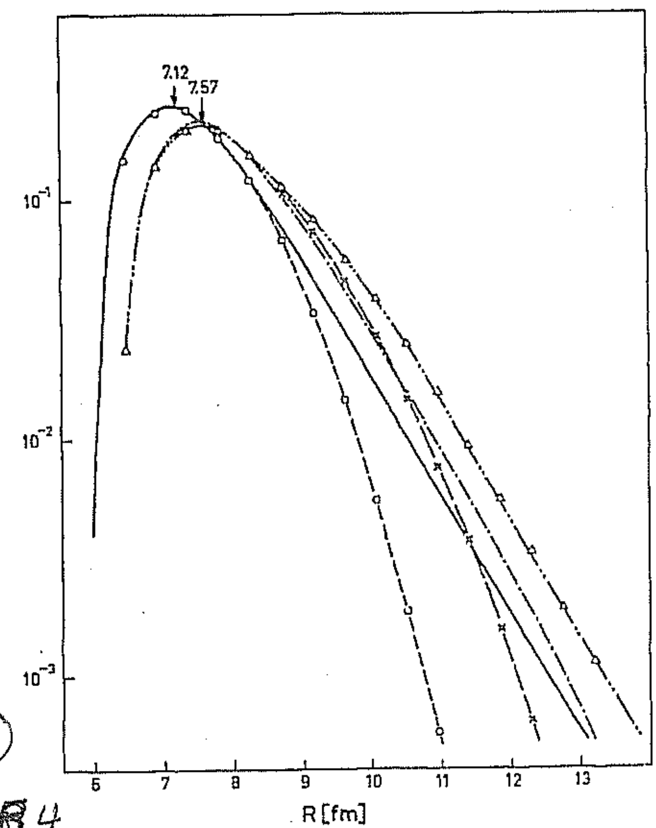


Fig. 7.B.3

~~Fig. 5.~~ Modified form factor for the transition to the ground state calculated in different spectroscopic models (pure shell-model configuration ———, shell model plus pairing residual interaction — — —, including ground state correlations —o—·—o—).



7.B.4

Fig. 7.B4

~~Fig. 6.~~ Asymptotic behaviour of the modified form factor for the ground state transition for oscillator plus Hankel wave functions (———), oscillator wave functions alone (—o— —o), using Woods-Saxon wave functions according to method I (—x—x—); (the same curve matched with a Hankel function ..— △..— △..—) and method II (—·—·—·).