ISTITUTO NAZIONALE DI FISICA NUCLEARE

Gruppo Collegato di Messina

INFN/BE-98/02
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Published by ***SIS-Pubblicazioni***
Laboratori Nazionali di Frascati

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ANGULAR CORRELATIONS AS A POWERFUL TOOL IN THE STUDY OF HEAVY–ION INDUCED PERIPHERAL REACTION MECHANISM

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Abstract

The $b - \alpha$ angular correlation analysis in $A(a,b)B(\alpha)C$ processes can be used as a powerful 'tool' to study the unbound states populated in transfer reactions leading to the same final state. So one obtains a complementary investigation of the intermediate nucleus $B \equiv \alpha + A$ studied by α-particle resonance scattering from A nucleus.

1 Sequential and two-body processes.

In Heavy-Ion Deep Inelastic Collisions (DIC) the mechanism governing both the transfer of most of the incident kinetic energy to the internal degrees of freedom of the intermediate system and the relative orbital angular momentum to the intrinsic angular momenta of the final reaction fragments has been extensively studied [1].

While the analysis of angular distributions in such reactions cannot give information on reaction mechanism because of the 'structurelessness' of the spectra, the particle-particle angular correlations , critically depending on the transition amplitudes, can be effective tools to study the reaction dynamics [2, 3].

2 Particle-Particle Angular Correlation.

Let us consider such a sequential process $A(a,b)B(c)C$ assuming that it proceeds through a given continuum state $(\epsilon^*_B, J_B\pi_B)$ in the nucleus B to a narrow definite state $(\epsilon^*_C, J_C\pi_C)$ in the final nucleus C [4].

ϵ^*_X is the excitation energy of the state having definite spin J_X and π_X in the nucleus X , while m_X is the z-component of $\vec{J}_X$. The (xX) system has relative radial coordinate $\vec{r}_x$, momentum $\vec{k}_x$, velocity $\vec{v}_x$ and energy ϵ_x. We define the spherical polar angles (ϑ_a, φ_a) of $\vec{k_a}$ and (ϑ_b, φ_b) of $\vec{k_b}$ in the $(A+a)$ centre-of-mass (c.m.) system, while $\vec{k_c}$ has polar angles (ϑ, φ) defined in the recoil centre-of-mass (r.c.m.) system (that is the rest frame of B) and described in a xyz-frame with the x-axis and z-axis parallel to the x-axis and z-axis of the c.m. frame.

In order for the $A(a,b)B(c)C$ reaction to be a sequential process, we require that the excitation energy ϵ_B^* of the intermediate system B formed in the first step of the three-body reaction be independent of the angles of the particle c and we also assume that in the $B \to c + C$ decay the nuclear interaction between b and B is negligible.

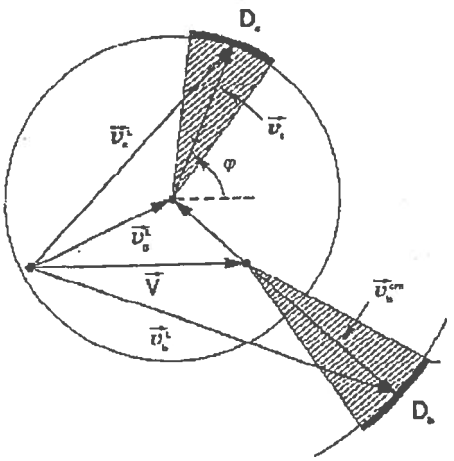

Figure 1: C.M. and r.c.m. Reference Frames.

Following the same symbolism of Ref. [4], we suppose for simplicity the nuclei A, a, b and c to have spin zero and b and c particles to be in the ground state.

Let us write the outgoing part of the stationary solution of the Schrödinger equation for the $A(a,b)B[\epsilon_B^*]$ reaction, corresponding to the partition $(bB; \epsilon_B^*, J_B m_B)$, associated with the unit flux in the partition $(aA; \epsilon_a)$, in the

asymptotic region ($\omega \equiv (\vartheta, \varphi)$) [5, 6]:

$$\Psi_b = f_{ba}(\epsilon_B^*; m_B, \omega_b)\phi_b\phi_B\mathcal{O}_b. \tag{1}$$

where the outgoing wave $\mathcal{O}_b$ is given by

$$\mathcal{O}_b = (v_b^{1/2}r_b)^{-1}\exp[i(k_br_b - \eta_b \ln 2k_br_b + \sigma_b^0)] \tag{2}$$

whilst η_b and σ_b^0 are the Coulomb field parameter and the S-wave Coulomb phase shift for the pair (bB), respectively; ϕ_b and ϕ_B are the normalized wave functions of the b and B nuclei, respectively; the f_{ba} amplitude , related to the differential cross-section for the $A(a,b)B$ reaction by means of

$$\frac{d\sigma}{d\omega_b} = \sum_{m_B} \mid f_{ba}(\epsilon_B^*; m_B, \omega_b) \mid^2, \tag{3}$$

determines the population of the substate $(\epsilon_B^*; J_Bm_B)$, in the nucleus B, prepared by detecting the particle b in the fixed direction of $\vec{k_b}$.

We suppose that, outside the range of any interaction with the nucleus b, the nucleus B decays into $c + C$ [4]. In the B nucleus rest frame, the outgoing part of the stationary solution of the Schrödinger equation corresponding to the decay from a substate $(\epsilon_B^*; J_Bm_B)$ to a given partition $(cC; \epsilon_C^*, J_Cm_C)$, in the asymptotic region can be written as [7]

$$\phi_B = F_{Bc}(\epsilon_B^*, \epsilon_C^*; m_B, m_C, \omega)\phi_c\phi_C\mathcal{O}_c. \tag{4}$$

where $\mathcal{O}_c$ is written in an analogous way like eq.(2), while ϕ_c and ϕ_C are the normalized wave functions for c and C particles, respectively; moreover, the decay amplitude F_{Bc} is defined by:

$$F_{Bc}(\epsilon_B^*, \epsilon_C^*; m_B, m_C, \omega) = -i\sum_\ell (-)^\ell < J_C\ell m_C, m_B - m_C \mid J_Bm_B > \cdot$$

$$\cdot\mathcal{S}_\ell(\epsilon_B^*, \epsilon_C^*)\ Y_\ell^{m_B - m_C}(\omega). \tag{5}$$

In the above equation ℓ is the relative orbital angular momentum of the pair (cC), $Y_\ell^m(\omega)$ are the spherical harmonics, $< J_C \ell m_C, m_B - m_C \mid J_B m_B >$ are the Clebsch-Gordan (CG) coefficients [9], $\mathcal{S}$ is the matrix which describes the $B \to c + C$ decay and the summation over ℓ is restricted to only parity-conserving values.

Now in order to evaluate, in the nucleus B rest frame, the outgoing part of the wave function which describes the sequential decay process $(aA; \epsilon_a) \to (bB; \epsilon_B^\star, J_B) \to (bcC; \epsilon_C^\star, J_C m_C)$ in the asymptotic region, we can use the above wave function ϕ_B thus getting, in the nucleus B rest frame

$$\Psi_{bc} = \mathcal{F}_{bc}(\epsilon_B^\star, \epsilon_C^\star; m_C, \omega_b, \omega) \mathcal{O}_b \mathcal{O}_c \phi_b \phi_c \phi_C \tag{6}$$

where

$$\mathcal{F}_{bc}(\epsilon_B^\star, \epsilon_C^\star; m_C, \omega_b, \omega) = \sum_{m_B} \mathcal{I} f_{ba}(\epsilon_B^\star, m_B, \omega_b) F_{Bc}(\epsilon_B^\star, \epsilon_C^\star; m_B, m_C, \omega). \tag{7}$$

Here $\mathcal{I}$ is the Jacobian related to the trasformation from the c.m. system to the r.c.m. system [6, 4] (see Fig. 1).

The number of the nuclei c emitted in the solid angle $d\omega$ along $\vec{k_c}$ in coincidence with the nuclei b emitted in the solid angle $d\omega_b$ along the fixed $\vec{k_b}$ is given by the flux j_{bc} in the partition $(bcC; \epsilon_C^\star, J_C m_C)$ through the surface elements of area $r_b^2 d\omega_b$ and $r_c^2 d\omega$, respectively, and defined by [5, 6]

$$j_{bc}(m_c; \omega_b, \omega) = v_b r_b^2 d\omega_b v_c r_c^2 d\omega \int d\zeta \mid \Psi_{bc} \mid^2 \tag{8}$$

the integration being taken over the totality of the internal variables.

If one recalls the unit flux in the partition $(aA; \epsilon_a)$, j_{bc} is exactly the double differential cross-section $d^2\sigma$ for the sequential reaction $(aA; \epsilon_a) \to (bB; \epsilon_B^\star, J_B) \to (bcC; \epsilon_C^\star, J_C m_C)$. So, by performing the integration in eq. (8),

with the functions ϕ_b, ϕ_c and ϕ_C normalized, the double differential cross-section for any final spin, from eqs. (6)-(8), is

$$\frac{d^2\sigma}{d\omega_b d\omega} = \sum_{m_C} \mid \sum_{m_B} F_{ba}(\epsilon_B^*; m_B, \omega_b) F_{Bc}(\epsilon_B^*, \epsilon_C^*; m_B, m_C, \omega) \mid^2 \tag{9}$$

where

$$F_{ba}(\epsilon_B^*; m_B, \omega_b) \equiv \mathcal{I} f_{ba}(\epsilon_B^*; m_B, \omega_b). \tag{10}$$

A quantity often considered in the $b-c$ coincidence measurements is the $b-c$ differential multiplicity defined by (see eqs. (3) and (9))

$$M(\omega) = \left(\frac{d^2\sigma}{d\omega_b d\omega}\right) \Big/ \left(\frac{d\sigma}{d\omega_b}\right), \tag{11}$$

that is the number of particles c per nucleus b and per unit solid angle in the moving frame of the decaying nucleus B, having sr^{-1} dimensions.

Let us take into account the same process in the particular case when $\omega_b = 0$ so that $m_B = 0$. With these assumptions, eq. (9) becomes:

$$\left(\frac{d^2\sigma}{d\omega_b d\omega}\right)_{\omega_b=0} = \mid F_{ba}(\epsilon_B^*) \mid^2 \sum_{m_C} \mid F_{Bc}(\epsilon_B^*, \epsilon_C^*; m_C, \omega) \mid^2 \tag{12}$$

from where one can deduce that the double differential cross-section can be factored in two terms, one related to the $A(a,b)B$ reaction, and the other one related to the $B \to c + C$ decay:

$$\left(\frac{d^2\sigma}{d\omega_b d\omega}\right)_{\omega_b=0} = \left(\frac{d\sigma}{d\omega_b}\right)_{\omega_b=0} \left(\frac{d\sigma}{d\omega}\right) \tag{13}$$

where

$$\left(\frac{d\sigma}{d\omega_b}\right)_{\omega_b=0} = \mid F_{ba}(\epsilon_B^*) \mid^2 \tag{14}$$

in which

$$F_{ba}(\epsilon_B^*) = -i(2\pi/k_a) \sum_{\ell_a \ell_b} S^{ba}_{\ell_a \ell_b} \sum_M (Y^M_{l_a}(\omega_a))^* \cdot \\ \cdot \langle I_B \ell_b 0, M \mid \ell_a M \rangle Y^M_{\ell_b}(0) \tag{15}$$

and

$$\frac{d\sigma}{d\omega} = \sum_{m_C} \mid F_{Bc}(\epsilon_B^*, \epsilon_C^*; m_C, \omega) \mid^2 \tag{16}$$

in which

$$F_{Bc}(\epsilon_B^*, \epsilon_C^*; m_C, \omega) = -i \sum_{\ell} (-1)^{\ell} S_{\ell}^{Bc} \cdot$$
$$\cdot < J_C \ell m_C, -m_C \mid J_B 0 > Y_{\ell}^{m_C}(\omega)^*. \tag{17}$$

Let us now consider the $A(\lambda, c)C$ elastic scattering process and let us analyze it by means of the **A**ckhiezer-**P**omeranchuk- **McI**ntyre-**B**lair model (**APMB**), where the particle scattering is treated by semi-classical boundary conditions.

The starting point is that the whole particle flow impinging on the nucleus is absorbed, while the other particles are scattered by the Coulomb potential exerted by the nucleus.

Let us write the cross-section for the elastic scattering

$$\frac{d\sigma}{d\omega} = \mid f(\theta) \mid^2 \tag{18}$$

where

$$f(\theta) = (2ik)^{-1} \sum_{\ell=0}^{\infty} (2\ell + 1)(\eta_{\ell} - 1) P_{\ell}(\cos\theta) \tag{19}$$

with $k = h/\mu v$ where μ is the reduced mass and v is the relative velocity, η_{ℓ} is the scattering matrix, $P_{\ell}(\cos\theta)$ are the Legendre polinomials.

We assume $\eta_{\ell} = 0$ and $\eta_{\ell} = e^{2i\sigma_{\ell}}$ for $\ell < \ell'$ and $\ell > \ell'$, respectively, while σ_{ℓ} is the Coulomb phase shift given by:

$$\sigma_{\ell} = \sigma_0 + \sum_{m=1}^{\ell} \tan^{-1}(\eta/m) \tag{20}$$

with

$$\eta = \frac{ZZ'e^2}{\hbar v}. \tag{21}$$

If one consider classical orbits, ℓ' can be expressed by:

$$\hbar^2\ell'(\ell'+1) = 2\mu R^2[E - (ZZ'e^2/R)] \tag{22}$$

where R is the sum of the radii of the involved nuclei, Z and Z' are their atomic numbers, E is the beam energy, and μ is the reduced mass.

To get the best fits at backward angles, we introduced a smooth cut-off for η_ℓ:

$$\eta_\ell = A_\ell e^{[2i(\sigma_\ell+\delta_\ell)]} \tag{23}$$

where δ_ℓ is a nuclear phase shift, while A_ℓ and δ_ℓ are given by [10]

$$A_\ell = \{1 + \exp[(\ell_A - \ell)/\Delta\ell_A]\}^{-1} \tag{24}$$

$$\delta_\ell = \delta_0\{1 + \exp[(\ell - \ell_\delta)/\Delta\ell_\delta]\}^{-1}. \tag{25}$$

The final differential cross-section expression becomes:

$$\frac{d\sigma}{d\omega} = (2k)^{-2} \mid \eta \sin^{-2}(\theta/2) \exp(2i\sigma_0) \exp[-i\eta \log(\sin^2(\theta/2))] + i\sum_{\ell=0}^{\infty} \exp(2i\sigma_\ell)(A_\ell \exp 2i\sigma_\ell - 1)P_\ell(\cos\theta) \mid^2 . \tag{26}$$

If the reaction cross-section is negligible with respect to the resonance energy, $A_\ell \simeq 1$ when ℓ is the resonance value; in this case, the Legendre polynomial coefficient amplitude reaches a maximum when $\delta_\ell = \pi/2$,and this is chosen as a resonance condition.

The corresponding level is attributed a spin $J = \ell$ and a parity $\pi = (-1)^\ell$. If we compare the expression of the differential cross-section with the general form obtained by partial wave analysis, δ_ℓ's represent the real part of partial wave phase shift and A_ℓ corresponds to $\exp(2i\delta_{\ell J})$, where the $\delta_{\ell J}$'s are the imaginary part of the partial wave phase shift.

Now we assume that the reaction $A(a,b)B$ proceeds via direct one-step transfer of a zero-spin particle λ, namely that the $A(a,b)B(c)C$ sequential process is represented by the sequence $a + A \to (b + \lambda) + A \to b + (A + \lambda) \to b + B \to b + c + C$. Then a visual comparison between eq.(13) and elastic scattering cross- section (18) shows the analogy between the $A(a,b)B(c)C$ sequential process and the $\lambda + A \to B \to c + C$ resonance scattering one as schematically depicted in Fig. 2 [8] by means of *Feynman-like* diagrams.

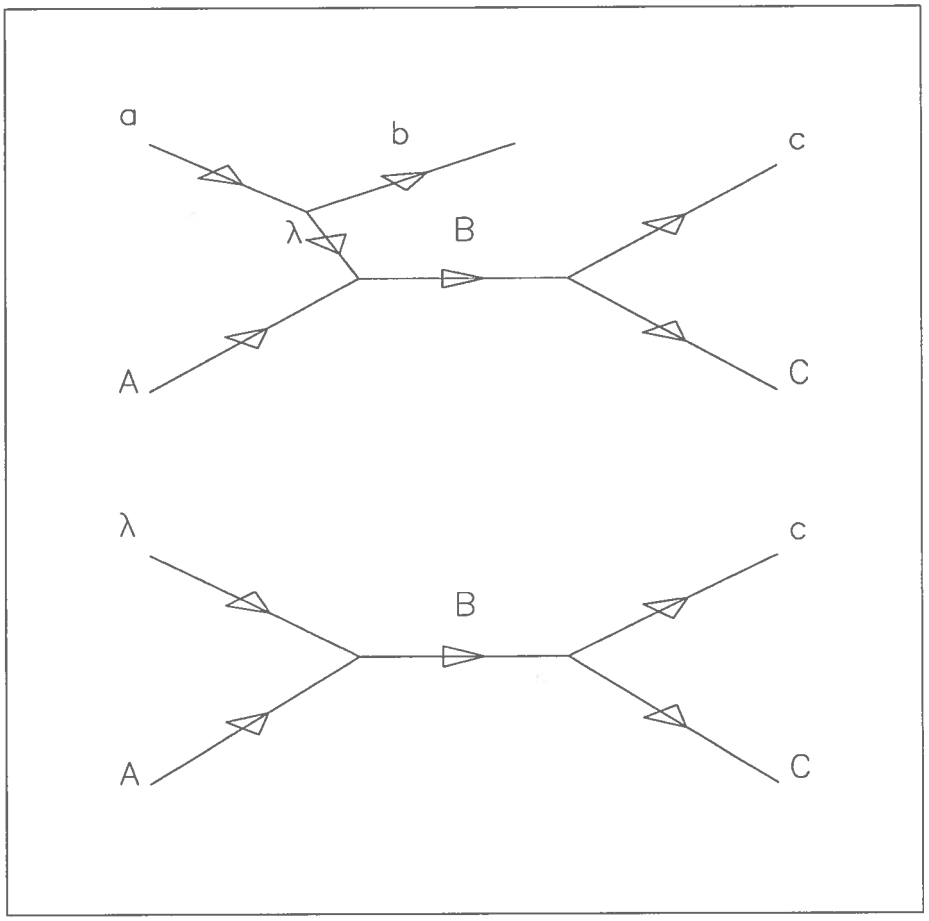

Figure 2: Diagramatic representation of a $a + A \to (b + \lambda) + A \to b + (A + \lambda) \to b + B \to b + c + C$ direct sequential process and the $\lambda + A \to B \to b + C$ resonance reaction.

Since our only assumption was that the $A(a,b)B$ reaction is a one-step transfer process, one can easily understand the decisive role played by $b-c$ angular correlation measurements in the study of the reaction mechanism. Different (a,b) reactions, leading to the same B nucleus, can be useful in the investigation of the role of the *spectator* b particle in more detail.

3 Analogy between sequential and two-body processes.

Instead of considering the coincidence cross-section by referring to specific energies of the nuclei B and C, we now want to get a theoretical expression of the b–c angular correlation for the three-body sequential process $A(a,b)B(c)C$ under the hypothesis that the B nucleus is highly excited to levels lying in the continuum region, within a certain, fixed energy interval; moreover and for simplicity let us consider only the case of a single J_B-value.

By averaging the double differential cross-section over an energy interval Δ centered at the energy $\epsilon_B^\star$ with a proper definition of the average of a function such that the final result of averaging should not depend in general on the exact form of the weighting function for the result to have physical sense, finally we get [4]

$$\langle \frac{d^2\sigma}{d\omega_b d\omega} \rangle \simeq \sum_{m_C} \mid \sum_{m_B} F_{ba}(m_B,\omega_b)\langle F_{Bc}(m_B,m_C,\omega)\rangle \mid^2 . \qquad (27)$$

Similarly to the one-level treatment, let us assume, also in the case of continuum, $m_B = 0$ and $\omega_b = 0$; then, from eq. (27):

$$\langle \left(\frac{d^2\sigma}{d\omega_b d\omega} \right)_{\omega_b=0} \rangle_{\Delta} \simeq \sum_{m_C} \mid F_{ba}(0^o)\langle F_{Bc}(m_C,\omega)\rangle \mid^2 \qquad (28)$$

one gets:

$$\langle\left(\frac{d^2\sigma}{d\omega_b d\omega}\right)_{\omega_b=0}\rangle_{\Delta} \simeq \mid F_{ba}(0^o) \mid^2 \sum_{m_C} \mid \langle F_{Bc}(m_C,\omega)\rangle \mid^2 . \tag{29}$$

From the previous equation, one can easily see that the double differential cross-section can be separated in two terms:

$$\left(\frac{d\sigma}{d\omega_b}\right)_{\omega_b=0} = \mid F_{ba}(0^o) \mid^2 \tag{30}$$

and

$$\left(\frac{d\sigma}{d\omega}\right) = \sum_{m_C} \mid \langle F_{Bc}(m_C,\omega)\rangle \mid^2 . \tag{31}$$

We want now to stress how, in this particular case, the differential multiplicity is exactly equal to the cross- section of the $B \to c + C$ decay:

$$\left(\frac{d\sigma}{d\omega}\right) = M(\omega) = \mid \sum_{\ell} \sum_{m_C} S_\ell(\epsilon_B^* + i\Delta) Y_\ell^{m_C}(\omega)^* \mid^2 \tag{32}$$

Since we expressed the reaction cross-section as a sum of an equilibrium term and of a non-equilibrium one, they will be factored as follows:

$$\left(\frac{d^2\sigma}{d\omega_b d\omega}\right)^E_{\omega_b=0} = \left(\frac{d\sigma}{d\omega_b}\right)^E_{\omega_b=0} \left(\frac{d\sigma}{d\omega}\right)^E \tag{33}$$

and

$$\left(\frac{d^2\sigma}{d\omega_b d\omega}\right)^{NE}_{\omega_b=0} = \left(\frac{d\sigma}{d\omega_b}\right)^{NE}_{\omega_b=0} \left(\frac{d\sigma}{d\omega}\right)^{NE} \tag{34}$$

where

$$\left(\frac{d\sigma}{d\omega_b}\right)^E_{\omega_b=0} = \left(\frac{d\sigma}{d\omega_b}\right)^{NE}_{\omega_b=0} = \mid F_{ba}(0^\circ) \mid^2 \tag{35}$$

while

$$\left(\frac{d\sigma}{d\omega}\right)^E = \sum_{m_C} \sum_{\ell} \omega_\ell(J_C) \frac{T_\ell}{G} \mid Y_\ell^{m_C}(\omega)^* \mid^2 \tag{36}$$

and

$$\left(\frac{d\sigma}{d\omega}\right)^{NE} = \sum_{m_C} \mid \sum_{\ell} S_\ell(\epsilon_B^* + i\Delta) Y_\ell^{m_C}(\omega)^* \mid^2 . \tag{37}$$

Let us now go back to the elastic scattering process, by averaging eq.(18) on a Δ energy interval, centered at the mean energy of the particles. The average of $f(\theta)$ can be written as:

$$\langle f(\theta)\rangle = (2ik)^{-1}\sum_{\ell=0}^{\infty}(2\ell+1)\langle(\eta_\ell-1)\rangle P_\ell(\cos\theta) \tag{38}$$

from which one gets:

$$\left(\frac{d\sigma}{d\omega}\right) = (2k)^{-2}\mid\sum_{\ell=0}^{\infty}(2\ell+1)\langle(\eta_\ell-1)\rangle P_\ell(\cos\theta)\mid^2 . \tag{39}$$

4 Using the DWBA formalism.

Let us consider now the same process by using the DWBA formalism: in this framework the transition amplitude $\mathcal{A}_{ba}$ of the sequential process

$$a+A\rightarrow(b+\tau)+A\rightarrow b+(\tau+A)\rightarrow b+B\rightarrow b+c+C$$

can be written as:

$$\mathcal{A}_{ba}(J_B m_B;\epsilon_B^*,\theta_b) = -[\hat{J}_B S_{N'0}(a,b)S_{NJ_B}(B,A)]^{1/2}\cdot$$
$$\cdot\sum_{L_b}T^{m_B}_{J_B L_B}(N,N')Y^{m_B}_{L_b}(\theta_b,0) \tag{40}$$

with

$$T^{m_B}_{J_B L_B}(N,N') = (-1)^{m_B}\left[\frac{4\pi}{\epsilon_a\epsilon_b}\frac{k_b}{k_a}\right]^{1/2}\sum_{L_a}i^{L_a-L_b-J_B}\cdot$$
$$\cdot\langle L_b J_B m_B,-m_B\mid L_a 0\rangle I^{J_B}_{L_a L_b}(N,N'). \tag{41}$$

$S_{NL}(\lambda,\mu)$ is the spectroscopic amplitude for λ nuclear state that decays in μ and τ with quantum numbers N and L, L_a and L_b are the orbital angular momenta for ingoing and outgoing channels, and $I^{J_B}_{L_a L_b}$ is the radial integral.

When $m_B = 0$ and $\omega_b = 0$, one gets:

$$\langle \left(\frac{d^2\sigma}{d\omega_b d\omega} \right)_{\omega_b=0} \rangle_\Delta = S_{N'0}(a,b) S_{NJ_B}(B,A) [\sigma_{J_B}(N,N',0)]_{DW} \langle \frac{dB_0}{d\omega} \rangle_\Delta \qquad (42)$$

where

$$\langle \frac{dB_0}{d\omega} \rangle_\Delta = \sum_{m_C} \mid \sum_{\ell} a_\ell(J_B, m_C; \epsilon_B^*, \Delta) Y_\ell^{m_C}(\omega) \mid^2 \qquad (43)$$

with

$$a_\ell(J_B, m_C; \epsilon_B^*, \Delta) = i(-1)^\ell \sum_{J_C} S_\ell^{Bc}(\epsilon_B^* + i\Delta) \langle J_C \ell, m_C, m_C \mid J_B 0 \rangle \qquad (44)$$

and

$$[\sigma_{J_B}(N,N',0)]_{DW} = (\widehat{J_B})^2 \mid \sum_{L_b} T^0_{J_B L_B}(N,N') \mid^2 \qquad (45)$$

is the differential cross-section at 0^o, for the $A(a,b)B$ reaction as expected by DWBA.

Let us now consider the differential cross-section averaged on a Δ interval centred at ϵ_B^* for the $A(a,b)B$ reaction with $J_A = 0$ proceeding through an isolated resonance $(\epsilon_B^*, J_B \pi_B)$, i.e.:

$$\langle \frac{d\sigma(\epsilon_B^*)}{d\omega} \rangle_\Delta = \frac{\pi}{k_a^2} \cdot \frac{[\Gamma_{Ba}(\epsilon_B^*)]_{J_B}}{\Gamma_B} \langle \frac{dB_0}{d\omega} \rangle_\Delta \qquad (46)$$

Γ_{Ba} and Γ_B being the partial and total widths, respectively, and $\langle dB_0/d\omega \rangle_\Delta$ is given by eq.(43).

As previously deduced, from comparison of eq.(29), (39), and (42), (46),one easily infers the analogy between the sequential process and the resonance scattering one.

the phase shift analysis of the resonances observed in $^{16}O(\alpha,\alpha)^{16}O_{g.s.}$ elastic scattering [11, 12, 13]. In the analysis of these states one has to support the thesis according which the four transferred nucleons behave like an α particle, so allowing a description which employs the cluster model.

Fig. 3 shows a comparison between the d-α coincidence spectrum for $^{16}O(^6Li,d)^{20}Ne(\alpha)^{16}O_{g.s.}$ at $E_{^6Li} = 35.3$ MeV at $\theta_d = 10°$ and $\theta_\alpha = 164° \div 170°$ in lab. and the excitation functions for $^{16}O(\alpha,\alpha_o)^{16}O_{g.s.}$ at $\theta_{C.M.} = 149.4°$ and $\theta_{C.M.} = 154°$. The subscript o on α means that only the α-decay to $^{16}O_{g.s.}$ is considered.

The correspondence among the highly excited levels in excitation functions and in α_o-decay spectrum, as well as the fair agreement - although qualitative - in their form, strongly supports the conjecture of an α-cluster structure of the more pronunced ^{20}Ne resonances in the considered energy range. This shows how the angular correlation analysis can be used, at least in the case of strongly excited levels, as a complementary technique to solve the ambiguities arising from the excitation functions analysis. It is clear that, in the study of the excitation function of the $^{16}O(\alpha,\alpha)^{16}O$ reaction and in the measurement of $d-\alpha$ coincidences related to the $^{16}O(^6Li,d)^{20}Ne(\alpha)^{16}O$ process, ^{20}Ne states will appear that have important reduced widths for α-decay to $^{16}O_{g.s.}$, due to the direct α-transfer in the $(^6Li,d)$ reaction.

One can look at the d-α angular correlation in $^{16}O(^6Li,d)^{20}Ne(\alpha)^{16}O$ as the angular distribution of α elastic scattering at different energies.

In ^6Li,^7Li α-cluster transfer-induced reactions the shape of the spectra and angular distributions at small angles and high excitation energy of residual nucleus, are mainly due to the momenta distributions of the clusters in ^6Li,

^7Li nuclei.

The angular distributions for high excited states do not depend on the angular momentum transferred to the nucleus: this circumstance does not allow one to deduce the state quantum numbers from angular distributions[14].

For these reasons an interrelationship is necessary between the information extracted from $b - \alpha$ angular correlation analysis and the one from α particle-induced resonance scattering: these studies can help nuclear spectroscopy in highly excited α level analysis, as well as providing new information on collective effects in light nuclei.

Acknowledgments

Authors are glad to thank Prof. A. Strazzeri for helpful contribution to these work and stimulating discussions on the subject.

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