Ref. 1 - 72-1066-1

A Pedestrian Approach To The Intermediate Structure
In Photonnuclear Reaction In Light Nuclei*

F. B. Malik

Department of Physics, Indiana University
Bloomington, Indiana 47401, U. S. A.

and

M. G. Mustafa[†]

Physics Division, Oak Ridge National Laboratory**
P. O. Box X, Oak Ridge , Tennessee 37830, U. S. A.

NOTICE

This report was prepared as an account of work sponsored by the United States Government. Neither the United States nor the United States Atomic Energy Commission, nor any of their employees, nor any of their contractors, subcontractors, or their employees, makes any warranty, express or implied, or assumes any legal liability or responsibility for the accuracy, completeness or usefulness of any information, apparatus, product or process disclosed, or represents that its use would not infringe privately owned rights.

MASTER

*Invited talk delivered by F. B. Malik at the V Symposium on the Structure of Low-Medium Mass Nuclei on October 27, 1972.

[†]On a leave of absence from Bangladesh Atomic Energy Centre.

**Operated by Union Carbide Corp. for the USAEC.

MASTER

DISTRIBUTION OF THIS DOCUMENT IS UNLIMITED

14

In talking about the photonuclear or any other types of a reaction theory in nuclear physics, it has become customary to classify them into two distant categories, viz the direct reaction (referred to as DR) and the compound process (denoted further as CP). The basic underlying picture of a DR is outlined in Fig. 1a and 1b. A photon of energy $h\nu$ being incident on a nucleus knocks out one of the nucleons, leaving the rest of the system essentially unperturbed. At best, the effect of the rest of the nucleus can be incorporated in terms of a smooth background of an optical potential distorting the outgoing nucleons. This is reminiscent of the dominating process in the photoemission in an atom and this approach was applied ~~by~~ to $C^{12}(p,\gamma)N^{13}$ reaction by Breit and Yost¹ in their attempt to explain the experiment of ref. 2. Since then, this model has been applied to the capture of nucleons in MeV region³ and recently⁴ to the emission of high energy nucleons from C^{12} and O^{16} . The excitation functions characterized by a direct emission exhibit broad resonances having widths of the order of a MeV in light nuclei.

The other extreme model widely used in understanding thermal neutrons (or neutrons in KeV region) is the concept of the compound nucleus, originally introduced by Bohr and Kalcker⁵ to understand the neutron capture cross section observed by Fermi and his collaborators⁶. According to the picture of a compound nucleus, photons entering a target nucleus excites the entire nucleus and all nucleons participate in this excitation. When a nucleon has accumulated enough momentum to escape, the compound nucleus emits this nucleon and in the process gets de-excited. The excitation functions, where this physical picture dominates, are characterized by resonances of narrow widths. The key difference between these two concepts is that in the

case of the DR only one or at the most a few nucleons participates and the rest of the nucleus remains a spectator, whereas in a CP, the whole nucleus participates. Moreover, in a DR, a nucleon is transferred directly from an initial to a final state without going through an intermediate state, whereas in a CP the formation of the intermediate compound nucleus plays a central role.

In actuality, a physical situation may occur which cannot be adequately idealized by either of these two extreme models and demands an intermediate approach blending salient features of both of them. Examples of such a situation is easy to visualize. In the case of photonucleon emission a fraction of the transition may be direct, whereas another fraction can go via an intermediate state to the continuum. In principle, it is not necessary to restrict the intermediate state to be a single one but there may be a number of them. This sharing of strength between a direct process and an intermediate process is even more likely, if there exists a relatively simple particle type of bound states embedded in the single nucleon continuum and the electromagnetic transition probabilities to this state is large. In that case, a fraction of the nucleons may be emitted directly and the remaining fraction may be ejected via an electromagnetic transition to this intermediate state.

A number of variation of this type of process can occur. First, there may be a bunch of intermediate states close to one another and an electromagnetic transition of a given multipolarity from the ground state to all of these states are allowed. By close, we mean that the mean separation of these states are of the order of their particle emission widths. In that case, the excitation

function will have a series of overlapping resonances and exhibit a succession of peaks on a broad bump. In case, the level spacing is shorter than the half widths of particle emission, we may not see any peak, but see only a broad bump in excitation functions, even though there may be a number of intermediate states under this bump.

A somewhat more sophisticated variant of the same thing may also occur. This was, first pointed out by Lane, Thomas and Wigner⁷. We may visualize that a very strong intermediate state is surrounded by a series of closely spaced states as shown on the left side of Fig. 2. In the first approximation, the selection rule does not permit any transition from the ground state to these surrounding states (let us call them satellite states or satellite doorway states). However if this state is well described in the first approximation only by a shell model type or a particle-hole or any type of models, there is a residual interaction between this strong state and its surrounding weak states. As a result of this interaction a fraction of this strong state is mixed to its surrounding states and a situation similar to the one shown on the right side of Fig. 2 can arise. This makes then a transition from the ground state to all these states permissible and nucleons then may be emitted from these states. The excitation function of such a process may again exhibit either a series of peaks on a background of a broad bump or a broad bump. The former will be the case if either the radiative width or the decay width or the combination of them associated with these satellite states is comparable to that of the strong state (primary intermediate (or doorway) states --PIS), and if it is not, the latter is the case.

Thus the key points associated with the occurrence of intermediate structures in excitation function are:

- a) The process is not entirely compound. This implies that a reasonable strength for a direct transition is available.
- b) Occurrence of one or many closely spaced intermediate states embedded in single nucleon emission continuum. A variant of this general process is the occurrence of one or more PIS state coupled to a series of satellite states.
- c) The relative strength of both the radiative widths to these states and the decay widths from them.
- d) The transition strength to a PIS states is shared or re-distributed between a PIS state and its satellites. The total nucleon emission probabilities is again shared between a direct transition and the transition via one or more of these intermediate states.

Thus attention must be paid to the normalization of the total wavefunction in every stage of the analysis, if one wants to obtain proper magnitude of the cross sections.

Once we have defined our physical picture, the next step is to find a suitable mathematical language to implement this. There are a variety of options starting from R-matrix theory to Green's function technology open, but I prefer to try the simplest one, even though this may not be very sophisticated. To this end, we start with the total Hamiltonian and proceed in a fashion analogous to Feshbach's approach^{8,9} but avoid the explicit use of projection operators and adopt a theory based on Trefftz's treatment¹⁰ of dielectronic recombination. Furthermore, we cast the theory specifically for the case of photonuclear reaction but its application in other cases like

fine structure in analogue resonances is straight forward.

The present day experimental situation in the photonuclear reaction in light nuclei has recently been reviewed by Firk¹¹. Considerable structure observed in the giant dipole excitation region led to the suggestion^{12,13} that these are clear cases of intermediate coupling where a part of the transition occurs through intermediate states. These states are generated by the coupling of a strong single particle or particle-hole state to a series of other weak and more complicated states of Boson-like character. This simple model using coupling constant and the decay width as free parameters was successful in correlating the observed structure in O^{16} and Si^{28} . Recently, Shakin and Wang have applied the similar idea to the structure in ^{16}O and using five parameters, namely the coupling constant, energy shifts and the decay widths of PIS and satellite states. They have been successful in generating enough structure in the excitation function. Rowe and Wang¹⁵, Goncharov et al.,¹⁶ and Kamimura et al.¹⁷ have also shown that the strong particle-hole 1-states in C^{12} can be successfully coupled to other nearby states and the strong radiative E1 transition strength to this state is, then, fragmented and shared by these nearby states via this coupling. However no attempt has been made to compute the actual excitation function.

II. THEORY

1. Equations

The total wavefunction for A nucleons in the initial or the final state can be expanded in a complete orthonormal set

7.

$$\psi = \sum_c \psi_c + \sum_b \psi_b \equiv \psi_I + \psi_{II}, \quad (1)$$

where ψ_c refer to all continuum open and close channels and ψ_b refer to all channel wavefunctions which are bound. Thus

$$(\psi_c, \psi_{c'}) = \int dE \delta(E_c - E_{c'}) \delta_{cc'},$$

$$(\psi_b, \psi_{b'}) = \delta_{bb'}, \quad (\psi_c, \psi_b) = 0. \quad (2)$$

In order to specify ψ_c and ψ_b , we are to specify the corresponding Hamiltonian. Since we shall be considering the photonuclear reaction of the type $^{16}\text{O}(\gamma, p)^{15}\text{N}$, it is appropriate to split the total Hamiltonian $H(A)$ as follows

$$H(A) = H_0(A-1) + T(i) + V(i, A-1), \quad (3)$$

where $H_0(A-1)$ is the Hamiltonian of the $(A-1)$ system, $T(i)$ is the kinetic energy operator of the i th particle which is being scattered and $V(i, A-1)$ is the interaction between the i th particle and $(A-1)$ nucleons. Let $\phi_c(A-1)$ and $\phi_b(A-1)$ be the continuum and bound state wave functions of H_0 satisfying the Schrödinger equations

$$H_0 \phi_\beta = E_\beta \phi_\beta ; \quad \beta = b, c \quad (4)$$

Then the functions ψ_b or ψ_c are constructed from ϕ_β 's by taking an antisymmetric product

$$\psi(A) = [\phi_\beta(A-1) \phi_j(i)]_a ; \beta, j = b, c \quad (5)$$

To define ϕ_j , we take the scalar product $(\phi_j, (H-E)\psi) = 0$, and get the set of equations

$$\begin{aligned} [E - E_j - T(\vec{r}_i) - V_{jj}(\vec{r}_i)] \phi_j(\vec{r}_i) \\ = \sum_{j' \neq j} V_{jj'}(\vec{r}_i) \phi_{j'}(\vec{r}_i) \end{aligned} \quad (6)$$

($j, j' = b, c; b = 1 \dots N, c = 1 \dots N$)

V_{jj} , are hermitian matrix elements

$$V_{jj'} = (\phi_j(A-1), V(i, A-1) \phi_{j'}(A-1)). \quad (7)$$

The form of (6) is obtained if the antisymmetrization is neglected. This coupled set of differential equations will have to be replaced by a different coupled set of integro-differential equations, if the Pauli Principle is incorporated, explicitly¹⁶.

Specializing now to those cases of ϕ_j 's having spherical symmetry, we may write

$$\phi_j(i) = \xi_j(\alpha, \hat{r}_i) r_i^{-1} \chi_j(r_i), \quad (8)$$

where α represents other quantum numbers necessary to define the state ϕ_j . Displaying explicitly indices b and c , we can rewrite (6) to emphasize these two different classes of solutions distinguished by their asymptotic behaviors.

$$\begin{aligned} [E - E_c - T_c(r) - V_{cc}(r)] \chi_c(r) = & \sum_{c' \neq c} V_{cc'}(r) \chi_{c'}(r) \\ & + \sum_b V_{cb}(r) \chi_b(r) \end{aligned} \quad (9a)$$

and

$$\begin{aligned} [E - E_b - T_b(r) - V_{bb}(r)] \chi_b(r) = & \sum_{b' \neq b} V_{bb'}(r) \chi_{b'}(r) \\ & + \sum_c V_{bc}(r) \chi_c(r), \end{aligned} \quad (9b)$$

$$\text{where } T_m = -\frac{\hbar^2}{2\mu} \left[\frac{d^2}{dr^2} - \ell_m(\ell_m+1)/r^2 \right]; \quad m = c, b.$$

We now discuss different approaches to the study of photonuclear reactions done to date within the context of these sets of coupled equations. For $^{16}\text{O}(\gamma, n)^{15}\text{O}$ type of reactions, initial bound states of ^{16}O are usually taken to be the solutions of (9b) with V_{bc} , i.e., the interaction of a bound state to a continuum state, to be zero. Further discussion relates to different types of approximation used for the final state wave function.

2. Approximations

A. Direct Reaction

Extreme case of direct reaction is obtained by setting V_{cb} , V_{bb} and $V_{bb'}$, all equal to zero. In addition, the distorted wave approximation assumes $V_{cc'} = 0$. This corresponds to the physical picture that the outgoing

particle is merely scattered by the average potential and there is no bound state in the continuum.

If, however, V_{cb} is not taken to be zero exactly but assumed to be small, following refs. 19, 20, one can show that an energy averaged contribution of these terms smoothing out resonance structures leads to a complex potential. This conclusion can also be drawn within the context of Feshbach's formalism. Thus in a direct reaction model for the photonuclear reaction, the final state single nucleon wave function is a solution of the optical equation

$$[E - E_c - T_c(r) - \bar{V}_{cc} - i\bar{W}_{cc}] \chi_c(r) = 0, \quad (10)$$

where $\bar{V}_{cc} + i\bar{W}_{cc} = V_{cc} +$ the complex potential originating from the coupling V_{bc} .

Actually, the first direct reaction model for the photonuclear reaction was used by Breit and Yost¹ to explain the $^{12}\text{C}(p,\gamma)^{13}\text{N}$ experiment of Hafstad and Tuve². Only partial success by them to explain this data lies in using only the real potential V_{cc} in obtaining $\chi_c(r)$ from (10). Since then, this model has been used in a number of cases. In particular, the photonuclear reaction in ^{16}O and ^{12}C above 30 MeV excitation energy has been successfully analyzed in refs. 4, 21. Even in the giant dipole region, this process yields a substantial contribution. This point is, of course, important for the application of the intermediate coupling model discussed later because if the contribution from the direct process is negligible, the process is compound and can not be considered to be intermediate.

B. Bound State Calculations

The work of Elliot and Flowers²² Gillet and Vinh Mau²³

and Brown and his collaborators²⁴ is of this type. The wave function

$\psi_b(A)$ is a bound one particle - one hole state, i.e., $\phi_b(A-1)$ is the hole state and $\phi_b(i)$ is the bound excited single particle state. For example, in ^{16}O , $\phi_b(A-1)$ is a $1p_{1/2}$ or $1p_{3/2}$ hole state and $\phi_b(i)$ is a $1d_{5/2}$, or $1d_{3/2}$, or $2s_{1/2}$ single particle state. Since no continuum state is present, the interactions

$$V_{cc} = 0 = V_{cc}',$$

$$V_{cb} = 0 = V_{bc}.$$

The relevant coupled equation is the following:

$$[E - E_b - T_b(r) - V_{bb}(r)] \chi_b(r) = \sum_{b' \neq b} V_{bb'}(r) \chi_{b'}(r).$$

Using an appropriate residual interaction, the final state wavefunction and energies are obtained after diagonalization in the one particle - one hole basis. These calculations showed a 2-4 MeV shift in the energies of the unperturbed one particle - one hole states and also showed that the dipole transition strength between these states and the ground state is concentrated in one or two transitions.

The study of the structure¹⁵⁻¹⁷ in giant dipole region of C^{12} falls also in this category. In essence, all these calculations calculate discrete lines, which constitutes only the first step towards calculating cross sections.

Locations of particle-hole states within the standard shell model scheme have also been carried out in Si^{28} . The E1 transition rates from the ground to these 1-states exceeds substantially the observed cross sections and fail

to account for the fine structure. This is, indeed, indicative that these states may be sharing a part of their strength with other states.

C. Continuum State Calculations

These calculations attempt to solve the set of equations (9a) by incorporating $V_{cc} + V_{cb}$ term in terms of a smooth optical potential and χ_c , are expanded in terms of 1 particle (p) - 1 hole (h) shell model type of states having bound state properties²⁵. The antisymmetrization has been incorporated and it is found that the antisymmetrization primarily affects the magnitude of the dipole transition but does not change the structure of the cross section. The main feature of the two major peaks at 22.3 and 24.3 MeV is well accounted for by these calculations. However, the observed structure could not be reproduced and the computed cross section remains 2 to 3 times larger than the observation. Incorporation of 2p-2h states in the expansion of χ_c , does not improve the situation adequately²⁶. A series of studies involving interaction between a number of continuum states has also been reported²⁷.

D. Intermediate Coupling Approximation

While the continuum state calculations focus primarily on the interaction between the continuum states and incorporate the interaction between a continuum and a bound state either as a smooth background in terms of optical potential or as a well separated resonance, the intermediate coupling model demands that the strong 1p - 1h states are coupled to a number of other bound states

by a residual interaction and thereby shares its strength with them and the E1 transition occurs in two steps: First the electromagnetic transition takes place from the ground state to these intermediate states in continuum. These then decay by particle emission very much in the same way as in the autoionization process in atomic physics.

Duke, Malik, and Firk (DMF)¹³ have described the giant dipole resonances in light nuclei by an intermediate coupling model which attributes their fine structure and widths to the coupling between single-fermion states and collective bosons associated with the quasi-stationary stage of the reaction. In their analysis, they assumed a linear coupling in momentum space between these fermion and boson states, and they have approximated these states by harmonic oscillator type wave functions. With such wave functions the diagonalization of the Hamiltonian has been achieved. It is natural that such wave functions will only produce discrete lines in the excitation function corresponding to different transition matrix elements which represent only the first phase of the intermediate coupling. The next phase involves the computation of the decay probability.

The final cross section will be discrete if the particle decay width is assumed to be a δ -function. However, this is obviously unphysical and DMF used a normalized Lorentzian width of Γ_s to compute the final cross section. This Γ_s is related to the lifetime of the particle-hole (or rather PHS) state in the absence of the coupling term.

In connection with our two sets of coupled differential equations (9), the DMF approximation consists of

- a) neglect $V_{cc'}$ ($c \neq c'$),
- b) use $V_{bb'} = g_\lambda f(r)$ where g_λ is some simple coupling constant.

The subscript λ identifies the different single particle excited states the combination of which gives the giant resonance cross section,

- c) take $\chi_b(r)$ as harmonic oscillator wave functions, and
- d) perform a diagonalization to account for $V_{bb'}$, and obtain a new set of basis $\bar{\chi}_b$ and express cross section in terms of the resonant amplitudes and decay widths of these $\bar{\chi}_b$.

Using this simple model the structure of giant dipole resonances in ^{16}O and ^{28}Si has been fitted using coupling constant and decay widths as parameters.

The purpose here is to investigate this intermediate coupling model within the context of a reaction theory and to relate energy shifts, widths, and cross sections to two body matrix elements. This will, then, throw more light on the resonant nature of structures in the giant dipole region.

III. Formal Solution of the Coupled Equations in the Intermediate Coupling Approximation

To solve the system of equations (9), we neglect coupling between different continuum channels $V_{cc'}$ ($c \neq c'$) in (9a). Although there

is no a priori justification in neglecting these terms, it is hoped that its effect can be partly incorporated by modifying the diagonal interaction. Continuum coupled channel calculations emphasizing these terms, already, indicate that they do not account for fine structures, although they influence the width of the resonances and may shift their energies. Their neglect, on the other hand, helps us to obtain a simple algebraic solution and to see clearly how the resonating part of the cross section influences the non-resonating background or direct component and vice versa.

Using $t_m(r) = d^2/dr^2 - \ell_m(\ell_m + 1)/r^2$, $k_c^2 = (2\mu/\hbar^2)(E - E_c)$, $k_b^2 = (2\mu/\hbar^2)(E_b - E)$, and $U_{mn} = (2\mu/\hbar^2) V_{mn}$, we may rewrite (9a) and (9b) in this approximation as

$$[t_c(r) + k_c^2 - U_{cc}(r)] \chi_c(r) = \sum_b U_{cb} \chi_b \quad (11a)$$

and

$$[t_b(r) - k_b^2 - U_{bb}(r)] \chi_b = \sum_{b' \neq b} U_{bb'} \chi_{b'}(r) + \sum_c U_{bc} \chi_c(r) \quad (11b)$$

Following Mott and Massey [24], the formal solution of (11a) is

$$\chi_c(r) = \chi_c^{(0)}(r) \delta_{cc'} + \int K(r, r') \sum_b U_{cb}(r') \chi_b(r') dr', \quad (12)$$

where

$$K(r, r') = (-k_c)^{-1} [\chi_c^{(0)}(r_<) \chi_c^{(1)}(r_>) + i \chi_c^{(0)}(r_<) \chi_c^{(0)}(r_>)] \quad (13)$$

Here $r_< (r_>)$ is the smaller (larger) of r and r' . $\chi_c^{(0)}$ and $\chi_c^{(1)}$ are, respectively, the regular and irregular solutions of the homogeneous part of (11a) (i.e., the right hand side of (11a) is set zero) having the following asymptotic forms:

$$\chi_c^{(0)} \xrightarrow{r \rightarrow \infty} \sin(k_c r - \eta_c \ln(2k_c r) - \pi \ell_c/2 + \lambda_c + \sigma_c)$$

and

$$\chi_c^{(1)} \xrightarrow{r \rightarrow \infty} \cos(k_c r - \eta_c \ln(2k_c r) - \pi \ell_c/2 + \lambda_c + \sigma_c)$$

where $\eta_c = Z_1 Z_2 e^2 / \hbar v$; Z_1 is the charge of the residual nucleus, Z_2 , the charge of the outgoing particle and v , the velocity of the particle in continuum. σ_c and λ_c are, respectively, Coulomb and non-Coulomb phase shifts.

Two important trivial points may be noted: (i) One can add or subtract any potential in defining the homogeneous part of (11a), e.g., the imaginary smooth potential simulating the non-resonant effect of U_{bc} can be included in defining homogeneous solution and (ii) formerly, continuum-continuum coupling terms $\sum_{c'} U_{cc'} \chi_{c'}$, could have been incorporated in (12) by replacing single sum over b by a double sum over b and c' , provided all matrix elements $U_{cc'} \chi_{c'} \rightarrow 0$ at least exponentially. We also note that solution (12) along with appropriate boundary conditions and (11) are equivalent to defining a scattering matrix.

Although, formally, (12) appears to be a solution, in principle it is nothing but an integral representation of (11), since χ_b 's are still unknown. Following Trefftz, we write a general solution of χ_b in terms of an orthonormal set

$$\chi_b(r) = \sum_n a_{nb} \chi_{nb}(r) + \int dq a_{qb} \chi_{qb}(r) \quad (14)$$

with $\int \chi_{nb}^2 dr = 1$, $\int \chi_{qb} \chi_{q'b} dr = \delta(q - q')$ and $\int \chi_{nb} \chi_{nb'} dr = 0$. Coefficients a_{nb} and a_{qb} can now be obtained from (11b), noting that χ_{nb} and χ_{qb} satisfy homogeneous part of these equations.

$$(k_{nb}^2 - k_b^2) a_{nb} = \sum_{b' \neq b} \int \chi_{nb} U_{bb'} \chi_{b'} dr + \sum_c \int \chi_{nb} U_{bc} \chi_c dr \quad (15)$$

$$-(q^2 + k_b^2) a_{qb} = \sum_{b' \neq b} \int \chi_{qb} U_{bb'} \chi_{b'} dr + \sum_c \int \chi_{qb} U_{bc} \chi_c dr. \quad (16)$$

a_{nb} in (15) has clearly a resonant behavior when k_b^2 is very close to k_{nb}^2 . Near such a resonance all χ_b may be approximated as (Note that this is not the case of isolated resonance but a series of successive resonances)

$$\chi_b(r) \approx a_{nb} \chi_{nb}(r) ; \quad b = 1 \dots N \quad (17)$$

Putting (12) and (17) in (15) we get

$$\begin{aligned} \sum_{b'} [(k_{b'}^2 - k_{nb}^2) \delta_{bb'} - \Delta k_{bb'}^2 + i \Gamma_{bb'}] a_{nb'} \\ = - \sum_c \int_0^\infty \chi_{nb}(r) U_{bc}(r) \chi_c^{(0)}(r) dr, \end{aligned} \quad (18)$$

where

$$\begin{aligned} \Delta k_{bb'}^2 = & - \int_0^\infty \chi_{nb}(r) U_{bb'}(r) \chi_{nb'}(r) dr \\ & + \sum_c k_c^{-1} \iint \chi_{nb}(r) U_{bc}(r) \chi_c^{(0)}(r_<) \chi_c^{(1)}(r_>) \\ & \times U_{cb'}(r') \chi_{nb'}(r') dr' dr \end{aligned} \quad (19)$$

and

$$\Gamma_{bb'} = - \sum_c k_c^{-1} \int_0^\infty \chi_{nb}(r) U_{bc}(r) \chi_c^{(0)}(r) \int_0^\infty \chi_c^{(0)}(r) U_{cb'}(r) \chi_{nb'}(r) dr \quad (20)$$

Here $\Delta k_{bb'}^2$ is the energy shift matrix. The first term on the right hand side of (19) is the energy shift due to the interactions between different bound states and the second term represents the shift due to the interaction of the bound states with the scattering states. The matrix $\Gamma_{bb'}$ is connected with the width of the resonance. Both $\Delta k_{bb'}^2$ and $\Gamma_{bb'}$ depend on the nature of two body interaction.

(12) along with (17) and (18) represent a complete set of equations defining the theory of intermediate structure. Clearly for all large a_{nb} , the wave function (12) has two distinct parts: One is $\chi_c^{(0)}$ corresponding to the scattering by a standard potential (or optical model), and the other connects a bound state to a bound state like structure embedded in the continuum which, finally, decays to the continuum. Thus intermediate structures are clear cut resonances but unlike the cases of isolated resonances, they can often be overlapping because if b and b' are far apart, the first term in

(19) and off diagonal Γ_{bb} , [20] may be neglected leading to well separated resonance like behavior.

IV. Cross Sections

For completeness we note that the above formalism is suitable for the study of intermediate structure in elastic and inelastic scatterings. For these cases, the scattering amplitude $f_{cc'}$, is defined by the asymptotic condition on χ_c as

$$\chi_c \xrightarrow{r \rightarrow \infty} \delta_{cc'} \sin \left(k_c r - \eta_c \ln 2k_c r - \frac{\pi \ell_c}{2} \right) + f_{cc'} \exp i(k_c r - \eta_c \ln 2k_c r - \pi \ell_c / 2) \quad (21)$$

with $\phi_b \xrightarrow{r \rightarrow \infty} 0$ (r^{-2}) and ϕ_b and $\phi_c \xrightarrow{r \rightarrow 0} 0$. Then the differential cross section for inelastic process ($c \neq c'$) is proportional to $|f_{cc'}|^2$

$$\begin{aligned} |f_{cc'}|^2 &= \left| \sum_b a_{nb} \int \chi_c^{(0)} U_{cb} \chi_{nb} dr \right|^2 \\ &= \sum_{bb'} a_{nb}^* a_{nb'} \int \chi_{nb} U_{bc} \chi_c^{(0)} dr \int \chi_c^{(0)} U_{cb'} \chi_{nb'} dr. \end{aligned} \quad (22)$$

For the elastic scattering f_{cc} is given by (22) plus a contribution coming from the potential scattering U_{cc} . If a_{nb} and $a_{nb'}$ are smooth functions, there will be no resonance like structure in the cross section but otherwise one may obtain a series of resonance like structures which will look like isolated resonances only if $\Gamma_{bb'}$, and $\Delta k_{bb'}^2$, (for $b \neq b'$) are zero.

The differential photonuclear reaction cross section is given by

$$\frac{d\sigma}{d\Omega} = (2\pi/\hbar) (\rho_f / \text{Photon Flux}) \sum |\langle \psi_f | H_{Int} | \psi_i \rangle|^2, \quad (23)$$

where H_{Int} is the interaction between photon and nuclei; ψ_f and ψ_i are eigenfunctions of H . Σ refers to an average over initial and sum over the final spin orientations. For cases like $^{16}\text{O}(\gamma, n)^{15}\text{O}$, the initial ground state is

far away from the continuum states and one can approximate $\psi_i \approx \sum_b \psi_b$ where ψ_b are constructed from shell model wave function solving (9b). This corresponds to incorporating correlation in the initial state. In the intermediate coupling model $\psi_f = \sum_c \psi_c + \sum_b \psi_b$ where ψ_c and ψ_b are constructed from (5) using the following approximations

$$\chi_c(r) \approx \chi_c^{(0)} + \sum_b a_{nb} \int K(r, r') U_{cb}(r') \chi_{nb}(r') dr' \quad (24)$$

and

$$\chi_b(r) \approx a_{nb} \chi_{nb}(r), \quad (25)$$

where $K(r, r')$ is given by (13), χ_{nb} are solutions of homogeneous equation defining χ_b , and a_{nb} are obtained by solving the set of coupled equations (18), exactly or in a suitable approximation.

V. Case of One Intermediate State

The matrix element for the transition from the ground state to a final state ψ_f via an isolated single intermediate state involves a simple expression for a_b (if there is no other intermediate state).

$$a_b = - \frac{\int \chi_b(r) U_{bf}(r) \chi_f^{(0)}(r) dr}{k^2 - (k_b^2 + \Delta k_{bb}^2) + i \Gamma_{bb}} \quad (26)$$

$$\equiv \frac{\gamma_{bf}}{E - (E_b + \Delta E_b) + i \Gamma_{bf}/2}, \quad (27)$$

where k_b in (18) is simply written as k and the index n in other quantities has been dropped for convenience. γ_{bf} is the reduced width and is related to the partial width by

$$\gamma_{bf} \equiv - \int \chi_b(r) V_{bf}(r) \chi_f^{(0)}(r) dr. \quad (28)$$

and

$$\begin{aligned} \Delta E_b &= (2/\hbar v) \iint \chi_b(r) V_{bf}(r) \chi_f^{(0)}(r_{<}) \chi_f^{(1)}(r_{>}) \\ &\quad \cdot V_{fb}(r') \chi_b(r') dr' dr \end{aligned} \quad (29)$$

This is the expression derived by Fano [25] and leads to the expression for the Breit-Wigner resonance [26]. This is evident because the contribution to the matrix element M in (23) due to the resonant part of the wave function is

$$M = \frac{\langle \psi_i | H_{\text{Int}} | \psi_b \rangle \gamma_{bf}}{E - (E_b + \Delta E_b) + i \Gamma_{bf}/2}. \quad (30)$$

The cross section is then proportional to

$$\frac{\Gamma_{ib} \Gamma_{bf}}{(E - (E_b + \Delta E_b))^2 + \Gamma_{bf}^2/4} \quad (31)$$

with

$$\Gamma_{ib} = |\langle \psi_i | H_{\text{Int}} | \psi_b \rangle|^2 \quad (32a)$$

and

$$\Gamma_{bf} = \frac{4}{\pi V} \left| \int \chi_b(r) V_{bf} \chi_f^{(o)}(r) dr \right|^2. \quad (32b)$$

This is Breit-Wigner resonance.

In principle however M contains also a term proportional to $\langle \psi_i | H_{\text{Int}} | \chi_c^{(o)} \rangle$ which is the contribution of a direct transition and there is an interference between these two terms in the expression for cross section. This is important for the intermediate coupling case.

VI. Case of Overlapping Resonances and Many Intermediate States

In this case, one is to diagonalize the following equation

$$\begin{aligned} \sum_{b'} [(k^2 - k_{b'}^2) \delta_{bb'} - \Delta k_{bb'}^2 + i\Gamma_{bb'}] a_{b'} \\ = - \sum_c \int \chi_b(r) U_{bc} \chi_c^{(o)}(r) dr \end{aligned}$$

or symbolically

$$\sum_{b'} [k^2 \delta_{bb'} - (W_{bb'} - i\Gamma_{bb'})] a_{b'} = \sum_c \chi_c. \quad (33)$$

That is one is to solve the matrix equation

$$(k^2 \cdot \underline{1} - \underline{W} + \underline{\Gamma}) \underline{a} = \underline{X} \quad (34)$$

This can simply be achieved by a similarity transformation which diagonalizes $\underline{W} - i\underline{\Gamma}$, i.e., to find a $\underline{t}$ such that

$$(\underline{W} - i \underline{\Gamma}) \underline{t} = \underline{t} \underline{\Lambda} \quad (35)$$

and if $t_{\alpha\beta}$ and $\Lambda_{\alpha\beta}$ are elements of $\underline{t}$ and $\underline{\Lambda}$ then we have

$$\sum_{\beta} (W_{\alpha\beta} - i\Gamma_{\alpha\beta}) t_{\beta\nu} = t_{\alpha\nu} \Lambda_{\nu}$$

with $\Lambda_{\nu} = \delta_{\nu\nu} (k_{\nu}^2 - i\gamma_{\nu})$, $\underline{W} - i\underline{\Gamma}$ is a symmetric matrix, and $\underline{t}$ is an orthogonal matrix, i.e., $\sum_{\alpha} t_{\alpha\nu} t_{\alpha\nu} = \delta_{\nu\nu}$, and $\sum_{\nu} t_{\alpha\nu} t_{\beta\nu} = \delta_{\alpha\beta}$ and hence $\underline{t}^{-1} = \underline{t}^T$ (its transpose). Since $\underline{W} - i\underline{\Gamma} = \underline{t} \underline{\Lambda} \underline{t}^T$, we get

$$(k^2 \cdot \underline{1} - \underline{W} + \underline{\Gamma})^{-1} = \underline{t} (k^2 \cdot \underline{1} - \underline{\Lambda})^{-1} \underline{t}^T$$

yielding

$$\underline{a} = \underline{t} (k^2 \cdot \underline{1} - \underline{\Lambda})^{-1} \underline{t}^T \underline{X} \quad (36)$$

or

$$a_b = \sum_{\nu b'} \frac{t_{b\nu} t_{b'\nu} x_{b'}}{(k^2 - k_{\nu}^2) + i\gamma_{\nu}}, \quad (37)$$

and it again shows resonance like behavior at an energy k_{ν}^2 which is obtained by diagonalizing $\underline{W}$ and a width γ_{ν} obtained by diagonalizing $\underline{\Gamma}$.

VII. Model of DMF

Instead of performing the complete transformation (35), one can do a partial diagonalization which is good if $U_{bc} \ll U_{bb}$, as it is visualized by Lane, Thomas, and Wigner [27]. In that case Γ_{bb} , defined in (20) can be neglected for all off diagonal elements and Δk_{bb}^2 , is given by the first term of (19).

In this particular case, one can first diagonalize (11b) and obtain a new basis $\bar{\chi}_{nb}$ and expand χ_b in terms of these $\bar{\chi}_{nb}$, i.e., (14) is replaced by the basis set

$$\chi_b(r) = \sum_n \bar{a}_{nb} \bar{\chi}_{nb} + \int dq \bar{a}_{qb} \bar{\chi}_{qb}(r) \quad (38)$$

For this basis set, $\Delta k_{bb'}^2$, is given by the second terms of (19) only with χ_{nb} replaced by $\bar{\chi}_{nb}$ (we call it $\bar{\Delta k}_{bb'}^2$), and $\Gamma_{bb'}$, is given by (20) with χ_{nb} replaced by $\bar{\chi}_{nb}$ (we call it $\bar{\Gamma}_{bb'}$). If, therefore, one of the U_{cb} , in (19) and (20) is small, these new $\bar{\Delta k}_{bb'}^2$ and $\bar{\Gamma}_{bb'}$, can be neglected. This is the model used by DMF under the assumption that $U_{bb'} = g_\lambda f(r)$, and Γ_{bf} of (32b) is taken as a parameter.

Since in this case $\bar{\Delta k}_{bb'}^2 = 0$ and $\bar{\Gamma}_{bb'} = 0$, the coupled equations (18) (defined by basis $\bar{\chi}_{nb}$) decouples and $\bar{a}_b$'s can be found easily. To illucidate this point we write (18) for the case of two amplitudes $\bar{a}_b$ and $\bar{a}_{b'}$:

$$\begin{pmatrix} k^2 - \bar{k}_b^2 - \bar{\Delta k}_{bb'}^2 + i\bar{\Gamma}_{bb'} & -\bar{\Delta k}_{bb'}^2 + i\bar{\Gamma}_{bb'} \\ -\bar{\Delta k}_{bb'}^2 + i\bar{\Gamma}_{bb'} & k^2 - \bar{k}_{b'}^2 - \bar{\Delta k}_{b'b}^2 + i\bar{\Gamma}_{b'b} \end{pmatrix} \begin{pmatrix} \bar{a}_b \\ \bar{a}_{b'} \end{pmatrix} = \begin{pmatrix} \bar{\gamma}_{bf} \\ \bar{\gamma}_{b'f} \end{pmatrix}$$

If $\bar{\Delta k}_{bb'}^2 = 0$ and $\bar{\Gamma}_{bb'} = 0$ then both $\bar{a}_b$ and $\bar{a}_{b'}$, are given by

$$\bar{a}_i = \frac{\bar{\gamma}_{if}}{k^2 - (\bar{k}_i^2 + \bar{\Delta k}_i^2) + i\bar{\Gamma}_{ii}} \quad ; \quad i = b, b' \quad (39)$$

Obviously, (39) is true for an arbitrary number of intermediate states.

Actually, one may note that for (39) to hold, it is not essential that $\bar{\Delta k}_{bb'}^2 = 0$ and $\bar{\Gamma}_{bb'} = 0$. They hold for the more general case of $\bar{\Delta k}_{bb'}^2$, and $\bar{\Gamma}_{bb'}$, being independent of either b or b' , e.g., large but constant or $\bar{\Gamma}_{bb'}/\bar{\Gamma}_b = f_b$.

Even in the approximation (39) generalized to contain an arbitrary number of resonances, the cross section is not just proportional to

$$\sum_b \frac{\bar{\Gamma}_{ib} \bar{\Gamma}_{bf}}{(E - (\bar{E}_b + \Delta \bar{E}_b))^2 + \bar{\Gamma}_{bb}^2/4} , \quad (40)$$

but contains an interference term between different a_b 's.

We note once more that the key difference between an intermediate coupling with overlapping resonances and a simple case of overlapping resonances leading to either a back ground imaginary potential and/or fluctuations is that U_{bb} , in (11b) is non-negligible. In the limit where (37) or (39) is valid the structures accounted for in the intermediate coupling model are actual resonances. However, if the density of intermediate states are very large in comparison to their width, these resonances can loose their own identities leading to fluctuations in cross sections.

This completes our formalism of the intermediate structure and overlapping resonances, and a microscopic justification and an analysis of assumptions involved in the DMF model. Obviously, the theory is a general one and not restricted to photonuclear case.

III. APPLICATION

We shall discuss in detail the situation in the photonuclear reaction in O^{16} within the context of an intermediate structure treatment and mention progresses made in case of Si^{28} and C^{12} .

A. O^{16}

An important clue suggesting that an intermediate structure type of situation may be accompanying a nuclear reaction is to estimate the direct transition probabilities and compare them with observation. If the direct transition probabilities form a significant fraction of the observed cross section, the entire process is not clearly a compound nuclear one in the classic sense and may admit an intermediate structure type of analysis.

The work of Buck and Hill²⁵ indicates already the presence of a reasonable direct component in the giant dipole energy region of O^{16} . A detailed analysis done in ref. 21 shows that the direct transition probabilities form a reasonable fraction of the excitation function. This is shown in Fig. 3. The strength of the direct transition probabilities is roughly 10% of the observed data.

For a further analysis, we discuss equations 11 through 16. Implications of these equations are

- 1) The interaction between bound states χ_{nb} has not been completely diagonalized. If a complete diagonalization of the Hamiltonian with respect to χ_{nb} is achieved, U_{bb} 's

on the right sides of (15) and (16) are zero and no further fragmentation is possible. In that case, (15) and (16) define the problem of overlapping resonances only. Mathematically, it is immaterial how the diagonal part of the interaction is defined. The important point is to include in off-diagonal matrix elements all interactions, which are not included in the diagonal part.

- 2) k_{nb}^2 in (15) refers to entire bound state spectrum. Thus a model Hamiltonian which describes low lying excitation spectrum should also be used to specify k_{nb}^2 lying in single particle continuum. One cannot consider states in giant dipole region with total disregard of the low lying spectrum.
- 3) Total normalization of the wavefunction is to be maintained. If a E1 transition rate from the ground state to a single unperturbed principal state k_{nb}^2 is strong, and very weak to its satellite states, the consideration of the coupling terms $U_{bb'}$ between this original PIS state and its satellite redistributes the strength of that transition. Consequently, the original strength of the transition is reduced and gets redistributed giving rise to broadenings.
- 4) The correlation in the ground state (or the initial state) is to be incorporated. As shown in ref. 4, the inclusion of the ground state correlation influences the magnitude of the cross section. In case of O^{16} , the magnitude is reduced.

- 5) Actual location of states depend on ΔE_{ob} , the energy shift and the shape depends on two widths, the radiative width Γ_{ib} and the decay width Γ_{bf} . This is evident from (31).
- 6) Since resonances corresponding to states of intermediate coupling are closely spaced, in principle the total width is not a simple sum of their individual width but one is to diagonalize (34) and find new widths.

In actual computation all these points have not yet been fully incorporated.

The model described in refs. 12 and 13, uses an arbitrary coupling parameter which is linear in momentum, to couple strong particle-hole states to nearly Boson-like states. The decay width associated with the PIS is taken to be of Lorentzian type and its width is used as a free parameter. The computed cross section is shown in Fig. 4. The interesting part is to note that because of the coupling, E1 transition can take place to a number of adjacent states and the whole width can take a structure observed in the $O^{16}(\gamma, n_0) O^{15}$ experiment.

In Fig. 5, we present an alternative fit to the broad resonance at 22.3 MeV. It clearly indicates that if the level spacing of satellite states is further shortened but the level width is nearly kept the same, the line shape can again be fitted and although there are a number of strong lines folded in the resonance, the fine structure is not pronounced.

Shakin and Wang¹⁴ adopt basically the same approach and idea and try to generate these boson states from particle - hole states.

These are then folded in a Lorentzian function to generate the final cross section. The calculated result is compared with experimental yield curve in Fig. 6. The absolute value is off by a factor of two. However if the integrated cross section is normalized to experimental data and then the decay width is used as a fitting parameter, the theoretical calculation fits the experimental data.

The most likely reason for Shakin and Wang's inability to account for the absolute magnitude of the cross section lies in their not considering points (2) and (3) discussed earlier. They do not consider ground state correlation and investigate the extent of compatibility of their states with low lying spectrum of O^{16} . This point is discussed below in a schematic fashion.

One of the best O^{16} theoretical spectra is due to Zucker, Buck and McGrory²⁸. O^{16} is considered by them as a C^{12} core plus four particles sharing $1p_{1/2}$, $2s_{1/2}$ and $1d_{5/2}$ shell model orbitals. El transition rates from their ground states to all possible 1- states in 10 to 25 MeV excitation energy generated by their wavefunction is shown in Fig. 7. Clearly, their wavefunction is not in a position to explain strong dipole transition. This is, of course, not at all surprising because the main El strength originates from the transition of a $1p_{3/2}$ particle to 2s-1d shell and ZBM wavefunctions do not have a $1p_{3/2}$ component. However a simple shell model wavefunction for $1p_{3/2}$ and 2s-1d states do not produce enough strength at correct energies. This is shown in Fig. 8a. As it has been pointed out by Brown and his collaborators²⁴ a particle - hole type of residual interaction is in a position to

place them in a proper position: This is shown in Fig. 8b . The main problem after this is

- a) to reduce the strength of particle hole states
- b) to generate additional fragmentations in the giant dipole excitation energy region keeping the basic structure of the ZBM wavefunction in tact.

Obviously, a good way to solve the problem is to open up C^{12} core and add a $lp_{3/2}$ component to the ZBM wavefunction and then perform the diagonalization. However, this is a very formidable task and will not be discussed here. A simple way to realize the effect of such a procedure is to weakcouple a $lp_{3/2}$ hole to the ZBM wavefunction. The unperturbed energies of $l -$ states obtained by weak coupling a $lp_{3/2}$ hole respectively, with a $ld_{5/2}$ and $2s_{1/2}$ particle states of the ZBM spectrum is shown in Fig. 9a and 9b , respectively. A 6.00 MeV spin orbit splitting between $lp_{1/2}$ and $lp_{3/2}$ has been used to generate these states. From the actual particle-hole calculation of Brown and others^{23,24}, it is known that the 18.31 MeV state of Fig. 9a is shifted by about 4 MeV to 22.3 MeV, and the rest of the states involving $ld_{5/2}$ particle orbital are shifted by the amounts mentioned in the Fig. 10. The energy shifts of coupling a $lp_{3/2}$ hole to $ld_{5/2}$ and $2s_{1/2}$ particle orbitals are given in Fig. 10. Using these energy shifts but generating particle - hole states from the ZBM we get the observed strengths shown in Fig. 10 . This figure also includes calculated dipole strength for the $l-$ states generated from $(lp_{1/2})^{-1} ld_{3/2}$ and $(lp_{3/2})^{-1} ld_{3/2}$. These yield additional strengths around 18 and 25 MeV.

Thus the expected dipole strength for a transition from the ground state

of O^{16} within the context of the ZBM model is the sum of Fig. 7 and Fig. 10. The most interesting point to observe is that dipole strengths have now correct magnitudes. The consideration of the correlation in the ground state configuration and a proper normalization of the total wavefunction are primarily responsible for this. This was not the case in the work of Shakin and Wang who, therefore, failed to get correct magnitudes.

However, it is to be emphasized that ours is only a simple schematic analysis of the problem. To complete this analysis, one would like to estimate properly the actual energy shifts in a particle - hole calculation based on the ZBM wavefunction and then either to fold the discrete dipole strength in a Lorentzian decay function as it has been done in refs. 13 and 14 or a more thorough analysis can calculate the decay widths and additional energy shifts themselves from a two nucleon potential using (19) and (20) or (37). We should also keep in mind that actual location of a state need not be at a maximum of a peak if Δk_{bb}^2 , in (19) $\neq 0$.

Another complexity in comparing a theoretical calculation with an experiment arises from the fact that experimental results in a photoemission involves a radiative transition of all multipoles, whereas theoretical results, discussed so far, consider only dipole transition. Experimental results on the study of angular distribution indicate an asymmetry about 90° which may arise from other multipole transitions. A simple theoretical study of the excitation function due to a M1 and E2 direct transition clearly show that the net contribution of such processes is only a few per cent of a E1 transition. This is shown in Fig. 3. Thus we do not expect a large E2 or M1 strength in this excitation energy. It is, however,

interesting to note that this slight admixture of E2 and M1 transition can reproduce the observed angular distribution and polarization data. This is shown in Fig.11 and Fig.12 .

A note of disparity hangs over the entire analysis of the photonuclear process. Because neutrons and protons have been treated on the same footing, i.e., an effective charge has been used for all neutrons. This may be indicative of a serious drawback of using spherical shell model states as basis. One should ponder legitimately about this effective charge and open one's mind to deformed basis. Such an attempt by Kluge²⁹ has already produced encouraging results. In particular, about 6, $K = 1$ and 3 $K = 0$ states can be generated between 19 and 26 MeV and the dipole strength to these states from the ground state is rather large. This is shown in Fig. 13

12 That the ground state correlation plays an important role in determining the correct magnitude of the dipole strength is also demonstrated in case of C^{12} . In Fig.14 , the dipole strength is calculated in a particle - hole approximation using no correlation in the ground state. Fig. 14 shows also transitions to the same particle - hole states but with ground state correlation incorporated by the open-shell Tamm-Dancoff method formulated by Rowe³⁰. Although the incorporation of this ground state correlation yields a reduced strength, it cannot account for the fine structure. To obtain further fragmentation Long and Rowe¹⁵ performed an intermediate coupling calculation using $(1s)^{-1}(1p)^9$ and $(1p)^7(2s\ 1d)^1$ configurations and Oak Ridge-Rochester shell model code³¹. This generated 41 $J = 1^-$ and $T = 1$ states and the fragmentation of the dipole strength is shown in Fig.14 . A similar calculation with a much more restrictive

space and yielding somewhat less fragmentation has also been performed by Goncharova and Yudin¹⁶. A comparison of the center part of Fig. 14 with its lower part demonstrates that the fine structure originates from a redistribution of lp-lh strengths among a series of satellite states by an intermediate coupling.

For a proper comparison of the observed total photonucleon cross section, these discrete dipole strengths are to be folded in particle decay widths. Since the particle decay width in C^{12} is a few MeV wide, much of these dipole strengths can be buried underneath without showing a great deal of fine structure. However, this results in broad bumps. Moreover, the problem of overlapping resonances must now be solved more accurately using (37).

In C^{12} , as in the case of O^{16} , the photonucleon emission of a neutron poses the intriguing question of "the effective charge". Once again a deformed orbital approach to generate intermediate structure may provide a solution to this question. The preliminary work of Nilsson, Sawicki, and Glendenning³² looks promising.

Si²⁸

Fig.15 indicates the experimental results³³ of the neutron yields from $Si^{28}(\gamma, n) Si^{27}$. The experiment seems to contain three well separated broad major resonances, the envelop of each of which contains small resonances. The spacing of these small or minor peaks is fairly uniform for a particular major resonance, which is analogous to the typical feature of phonon broadened electronic transition in solids due to low energy vibration.

The particle - hole calculations³⁴ yields three 1- states in the 17-21 MeV excitation energy region but fails in magnitude and cannot explain these fine structures.

Although Si^{28} is a complicated nucleus, it is reasonable to visualize that a large number of collective states of Boson like characters is present at such a high excitation energy. In fact a simple estimate based on a classical liquid drop model indicates a level spacing of 30 to 80 keV between these states at 17-21 MeV. In that case, these strong particle - hole states can be coupled to these Boson-like states and this generates a situation which is very similar to the phonon-broadened electronic transition in solids. The intermediate coupling model used in refs. 12 and 13 is particularly suitable for this. The computed excitation function is shown in Fig. 16, and can account for the magnitude and the broadening of observed resonances.

IV. CONCLUDING REMARK

The emerging picture associated with excitation functions in the photo-nuclear reactions indicates that these comprise of a direct emission and a transition via a series of closely lying intermediate states. Thus these are clear cases of overlapping resonances where each resonance has not yet completely lost its individual identity to the extent that they can be termed as fluctuations. On the other hand, we may learn more about a sophisticated reaction mechanism by trying to understand in greater depth some of basic parameters like particle width from a two body potential.

This analysis also points towards a systematic change in the character of the observed structure across the periodic table. Whereas the large nucleon emission width in light nuclei can smear out resonances associated with complicated

states, as in C^{12} , the situation changes with heavier nuclei because of the systematic variation of two basic parameters with increasing mass number A . First the single particle decay width decreases rapidly with increasing A and secondly the density of collective states of Boson-like character increases with A . Because of this and of the finite experimental resolution the fine structure may no longer be observed beyond $A \sim 40$ or 50 and instead one would see a single large bump.

Clearly, this intermediate structure is not restricted to the photo-nuclear reaction only. In case of the (p,n) reactions through analogue states in light and medium nuclei, the similar physical situation persists. The strong analogue state can be fragmented by a coupling to its satellite state. Apart from the usual nuclear residual interaction, the Coulomb interaction couples such analogue states with other nearby complex particle or collective states.

REFERENCES

1. G. Breit and F. L. Yost, Phys. Rev. 48, 203 (1935).
2. L. R. Hafstad and M. A. Tuve, Phys. Rev. 47, 506 (1935); C. C. Lautrisen and H. R. Crane, Private communication to G. Breit and F. L. Yost (1935).
3. See e.g., G. A. Bartholomew, Ann. Rev. Nucl. Sc. 11, 259 (1961); or R. R. Christy and I. Duck, Nucl. Phys. 24, 89 (1961).
4. M. G. Mustafa and F. B. Malik, Phys. Rev. C1, 753 (1970).
5. N. Bohr and F. Kalckar, Kgl. Danske Videnskab. Selskab. Mat-fys. Medd., 14, No. 10 (1937).
6. E. Amaldi, O. D'Agostino, E. Fermi, B. Pontecorvo, F. Rosetti, and E. Segre, Proc. Roy. Soc. (London) A149, 387 (1935).

7. A. M. Lane, R. G. Thomas and E. P. Wigner, Phys. Rev. 98, 693 (1955).
8. H. Feshbach, Ann. Phys. (N.Y.) 5, 357 (1958).
9. H. Feshbach, Ann. Phys. (N.Y.) 19, 287 (1962).
10. E. E. Treffitz, Z. Astrophysik, 65, 299 (1967).
11. F. W. K. Firk, Ann. Rev. Nucl. Sc. 20, 39 (1970).
12. C. B. Duke and F. B. Malik, Proc. Int'l. Conf. on Nucl. Phys. held in Gatlinburg, Tennessee (1966).
13. C. B. Duke, F. B. Malik, and F. W. K. Firk, Phys. Rev. 157, 879 (1967).
14. C. M. Shakin, and W. L. Wang, Phys. Rev. Lett. 26, 902 (1971),
Phys. Rev. C5, 1898 (1972).
15. D. J. Rowe and S. S. M. Wong, Phys. Lett. 30B, 147, 150 (1969).
16. N. G. Goncharova and N. P. Yudin, Phys. Lett. 29B, 272 (1969).
17. K. Kamimura, K. Ikeda, and A. Arima, Nucl. Phys. A95, 129 (1967).
18. M. G. Mustafa, doctoral thesis, Yale University, 1970, (unpublished).
19. N. Austern, Ann. Phys. (N.Y.) 45, 113 (1967).
20. F. M. Nicolau, Z. Naturforsch. 16a, 603 (1961).
21. M. G. Mustafa and F. B. Malik, Bull. Am. Phys. Soc. 14, 607 (1969) and
Phys. Rev. C2, 2068 (1970).
22. J. P. Elliot and B. H. Flower, Proc. Roy. Soc. (London) A242, 57 (1957).
23. V. Gillet and N. Vinh Mau, Nucl. Phys. 54, 321 (1964).
24. G. E. Brown and M. Bolsterli, Phys. Rev. Lett. 3, 472 (1959); G. E. Brown,
L. Castillejo, and J. A. Evans, Nucl. Phys. 22, 1 (1961).
25. C. Mahaux and H. A. Weidenmuller, Phys. Lett 19, 408 (1965); B. Buck and
A. D. Hill, Nucl. Phys. A95, 271 (1967); J. D. Perez and W. M. MacDonald,
Phys. Rev. 182, 1066 (1969); W. P. Beres and W. M. MacDonald, Nucl. Phys.

- A91, 529 (1967); J. Raynal, M. A. Melkanoff, and T. Sawada, Nucl. Phys. A101, 369 (1967); A. M. Saruis and M. Marangoni, Nucl. Phys. A132, 433 (1969).
26. V. Gillet, M. A. Malkanoff, and J. Raynal, Nucl. Phys. A97, 631 (1967).
27. C. Mahaux and H. A. Weidenmuller, Shell Model Approach to Nuclear Reactions (North-Holland Publishing Co., Amsterdam, 1969); W. Glockle, J. Hufner, and H. A. Weidenmuller, Nucl. Phys. A90, 481 (1967).
28. A. P. Zuker, B. Buck, and J. B. McGrory, Phys. Rev. Lett. 21, 39 (1968); Brookhaven National Lab Report, BNL 14085 (1968).
29. G. Kluge, Z. Physik, 197, 288 (1966).
30. D. J. Rowe, Rev. Mod. Phys. 40, 153 (1968); Nucl. Phys. A107, 99 (1968).
31. J. B. French, E. C. Halbert, J. B. McGrory and S. S. M. Wong, Adv. in Nucl. Phys. Vol. III eds. M. Baranger and E. Vogt (Plenum Press, N.Y. 1969).
32. S. G. Nilsson, J. Sawicki, and N. K. Glendenning, Nucl. Phys. 33, 239 (1962).
33. F. W. K. Firk, Bull. Am. Phys. Soc. 11, 367 (1966).
34. J. B. Seaborn and J. M. Eisenberg, Nucl. Phys. 63, 496 (1965).
35. F. W. K. Firk and K. H. Lokan, Phys. Rev. Lett. 8, 321 (1962) and F. W. K. Firk, Nucl. Phys. 52, 437 (1964).
36. J. E. E. Baglin and M. N. Thompson, Nucl. Phys. A138, 73 (1969).
37. G. W. Cole Jr., F. W. K. Firk, and T. W. Phillips, Phys. Letters 30B, 91 (1969).
38. G. W. Cole, Jr., Ph.D. Thesis, Yale University, 1970 (unpublished).

FIGURE CAPTIONS

Figure 1a:

A schematic drawing of the physical situation underlying a direct process (top picture), and a compound nuclear process (center picture), and an intermediate process (bottom picture). In a direct reaction a photon essentially knocks out a single nucleon without affecting other nucleons appreciably. In a compound nuclear process, a photon interacts with a nucleus as a whole and this nucleus remains in an excited state until a nucleon has gathered enough momentum to escape. In an intermediate process a photon can interact with a number of nucleons. A part of the emitted neutron flux is by a direct process and another part is emitted by excited states of the nucleus.

Figure 1b:

Columns marked Direct, Compound and Intermediate refer, respectively, to first, second and third rows of Fig. 1a in terms of nuclear levels.

Figure 2:

A schematic representation of an intermediate reaction process. The first column represents unperturbed states generated by an approximate model Hamiltonian. Transitions from the ground state to a few states marked by long bars are allowed. These states then decay by a particle emission yielding an excitation function shown in the second column. If the model Hamiltonian used to generate

states shown in the first column is not an adequate approximation to the actual Hamiltonian, the states shown in column one, can get mixed by a residual interaction and a transition from the ground state to nearby satellite states are allowed. The excitation function may look like the one sketched in the last column.

Figure 3:

Calculated E1, E2, and M1 integrated cross sections in a direct reaction model for the reaction $O^{16}(\gamma, p)N^{15}$ as a function of γ -ray energy E_γ . σ_{E1} , σ_{E2} , and σ_{M1} , are, respectively, the E1, E2 and M1 integrated cross sections. σ_T is the sum of these three. (Taken from ref. 21).

Figure 4:

Heavy solid vertical lines are computed intensities of the spectral lines in a photonuclear process in O^{16} calculated using the model discussed in ref. 13. The dashed lines are expected cross section if the principal state has a width Γ_s . The light solid line represents the experimental points in the $O^{16}(\gamma, n)O^{15}$ reaction (ref. 35) obtained by Firk et al. (Taken from ref. 13.)

Figure 5:

The same as in Fig. 4 but for a different set of parameters. This demonstrates that if the particle emission width is large in comparison to level spacings, the underlying states may be difficult to detect.

Figure 6:

Experimental cross section σ for the reaction $O^{16}(\gamma, n_0) O^{15}$ (dots) are compared to theoretical calculation of ref. 14. The dashed curve is the actual calculation using the decay width as a parameter. The solid line is the theoretical cross section normalized to yield total integrated experimental value.

Figure 7:

El transition rates from the ground to 1- states in the giant dipole region of O^{16} calculated from the ZBM wavefunctions (ref. 28).

Figure 8a:

El transition rates from the ground state of O^{16} to 1- states in a single particle elementary shell model where a 1p nucleon is transferred to the 2s-1d shell.

Figure 8b:

El transition rates from the ground state of O^{16} to 1- particle - hole states generated by coupling a 1p nucleon hole state of an elementary shell model to the 2s-1d particle states.

Figure 9:

The upper half of this figure represents the computed (vertical bars) El dipole strength from the ground state of O^{16} to 1- states generated by weakcoupling a $1p_{3/2}$ hole to $(2s_{1/2})$ and $(1d_{5/2})$ states of ZBM (ref. 28). The actual energy shifts ΔE calculated using a δ -function force discussed in ref. 14 are 3.89, 0.21, 0.58 and 0.15 MeV, respectively, for states at 18.31, 22.25, 23.34 and

24.00 MeV in Fig. 10a and are 1.07, 0.20, 0.15, and 0.11 MeV for states at 18.66, 22.10, 24.68, and 25.20 MeV, respectively, of Fig. 10b. Vertical bars in the lower half of this figure represents the same E1 transitions if an average shift of 1.0 MeV is assumed for all except the 18.31 MeV state of Fig. 10a and an average shift of 0.5 MeV is assumed for all except 18.66 state of Fig. 10b. Solid curve in each figure represents experimental results (ref. 35).

9a

Figure 10a:

Unperturbed energies of 1- states generated by coupling a ($1p_{3/2}$) hole to a ($1d_{5/2}$) particle state of the ZBM (ref. 28). A 6 MeV spin-orbit splitting between the $1p_{3/2}$ and the $1p_{1/2}$ orbitals has been used.

9b

Figure 10b:

The same as in Fig. 10a for the coupling of a ($1p_{3/2}$) hole to a ($2s_{1/2}$) particle state.

Figure 11:

The measured angular distributions in the $O^{16}(\gamma, p_O)N^{15}$ reaction normalized to unity at 90° (c.m.) from ref. 36, marked as vertical bars are compared with theoretical results (solid curves) of direct emission of ref. 21.

Figure 12:

In the upper part of the figure, observed $O^{16}(\gamma, n_O)O^{15}$ polarization cross section at 90° from refs. 37 and 38 are plotted as vertical bars and compared to the theoretical prediction of a direct emission used in ref. 21. The solid

line in the upper half represents calculated results using a single effective charge of $\pm 0.5e$ for E1 and E2 transitions. Triangles are theoretical computations using neutron effective charges of $(-Z/A)$ and Z/A^2 for E1 and E2 respectively. The curve in the lower half of the figure is the theoretical proton polarization at 90° in the $O^{16}(\gamma, p_0)N^{15}$ reaction using a single proton effective charge of $+0.5e$ for E1 and E2 transitions.

Figure 13:

The dipole strength from the ground state rotational band to rotational states in the giant dipole region (taken from ref. 29).

Figure 14:

The uppermost part represents the theoretical calculation of dipole strength from the ground state of C^{12} to 1- states generated by a conventional 1 particle - 1 hole calculation. The center picture represents the dipole strengths to 1- states generated by an open shell Tamm-Dancoff calculation suggested in ref. 15. The lowermost figure represents E1 rates from the ground state of C^{12} to all 1- states generated by an intermediate coupling shell model calculation incorporating $(1s)^{-1}(1p)^9$ and $(1p)^7(2s1d)^1$ configuration space (taken from ref. 15). Dashed curve represents experimental cross section.

Figure 15:

The experimental neutron yield obtained for the reaction

$\text{Si}^{28}(\gamma, n) \text{Si}^{27}$ by Firk (unpublished). The yield below 19 MeV may be artificially enhanced by the inclusion of non-ground state transition. (Taken from ref. 13.)

Figure 16:

The vertical lines are E1 transition probabilities from the ground state of Si^{28} to a series of 1- states generated by coupling three strong particle-hole states to its satellites by a linear coupling. (Taken from ref. 30.) The lightly dashed line is inserted as a visual aid to locate the envelope of the sharp line-spectra. Parameters used are noted in the figure.

INITIAL

INTERMEDIATE

FINAL

NONE

SELECTION
RULE
NEGATIVE

NEUTRON

+

+ n

DIRECT

COMPOUND

INTERMEDIATE

Fig 16.

NO RESIDUAL
INTERACTION

STATES

EXCITATION
FUNCTION

WITH RESIDUAL
INTERACTION

STATES

EXCITATION
FUNCTION

Fig.2

Fig. 3

Fig. 4

Fig. 5

Fig. 6

Fig. 7

Fig. 8

(1p 3/2, 1d 5/2)

(a)

(1p 3/2, 2s 1/2)

(b)

Fig. 9

Fig. 10

Fig. 11

Fig. 12

Fig. 13

Fig. 14

Fig. 15

Fig. 16