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THE IMPORTANCE OF NUCLEON-NUCLEON INTERACTIONS IN NUCLEAR REACTIONS

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THE IMPORTANCE OF NUCLEON-NUCLEON INTERACTIONS IN NUCLEAR REACTIONS

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A brief history is traced of the ideas that brought us to the present understanding of the processes which lead to the formation of a compound nucleus in a state of statistical equilibrium.

In 1936 Bohr, on the basis of the results of neutron-induced processes studied by Fermi *et al* (1934, 1935a,b), Szilard (1935) and Frisch *et al* (1936a, 1936b), made the hypothesis that *"a collision between a high-speed neutron and an heavy nucleus will in the first place result in the formation of a compound nucleus of remarkable stability"*. He also suggested that *"the possible later breaking up of this intermediate system by the ejection of a material particle, or its passing with emission of radiation to a final state, must in fact be considered as separate competing processes which have no immediate connexion with the first stage of the encounter"* (what is now known as the Bohr independence hypothesis), and that *"the essential feature of nuclear reactions, whether incited by collision or radiation, may be said to be a free competition between all the different possible liberation of material particles and of radiative transitions, which can take place from the semi-stable intermediate state of the compound system"*. These hypotheses, according to Bohr, were not restricted to processes induced by particles with an energy of a few MeV, but *"even if we could experiment with neutrons or protons of energies of more than a hundred million volts, we should still expect that the excess energy of such particles, when they penetrate into a nucleus of not too small mass, would in the first place be divided among the nuclear particles with the result that a liberation of any of these would necessitate a subsequent energy concentration. Instead of the ordinary course of nuclear reactions we may, however, in such cases expect that in general not one but several charged or uncharged particles will eventually leave the nucleus as a result of the encounter"* .

Bohr did not attempt to give a precise dynamical description of the formation of the compound nucleus. He simply stated: *"If, for example, we consider an encounter between a high-speed neutron and a nucleus, it is obviously not permissible to compare the process to a simple deflection of the path of the neutron in the inner nuclear field, possibly combined with a collision with a separate nuclear particle, resulting in the ejection of the latter. On the contrary, we must realise that the excess energy of the incident neutron will be rapidly divided among all the nuclear particles with the result that for some time afterwards no single particle will possess sufficient kinetic energy* *to leave the nucleus. The possible subsequent liberation of a proton or an α-particle or even the escape of a neutron from the intermediate compound system will therefore imply a complicated process in which the energy happens to be again concentrated on some particle at the surface of the nucleus .... In the atom and in the nucleus we have indeed to do with two extreme cases of mechanical many-body problems for which a procedure of approximation resting on combination of one-body problems, so effective in the former case, loses any validity in the latter where we, from the very beginning, have to do with essential collective aspects of the interplay between the constituent particles"* .

The experimental data that suggested the CN hypothesis were the neutron resonances that occur when the reaction goes through a single state in the compound nucleus. The cross section of such reactions, as a function of energy, shows a pronounced resonance structure that may be reproduced by the Breit - Wigner (1936) theory that appeared shortly after the Bohr's paper and greatly contributed to the success of the CN theory. A typical example of resonances is given in Fig. 1 where are shown those excited in the interaction of low energy neutrons on ^{238}U (Firk *et al* 1960). Subsequent development of the CN concept were the statistical theory of nuclear reactions (Weisskopf-Ewing 1940 and Hauser and Feshbach 1952) and the theory of statistical fluctuations of nuclear cross sections that occur when the average width Γ of the CN levels is much greater than their average spacing and are observable when the cross-sections are measured with an energy resolution smaller than the width, $\Delta E \leq \Gamma$. Fluctuations occur as a consequence of the random character of the transition amplitudes for a CN reaction (Ericson 1960, 1963a,b and Brink *et al* 1963, 1964a,b).

However, even before the appearing of the CN hypothesis, Oppenheimer and Phillips (1935) had shown that (d,p) reactions at very low energies occur by direct transfer of a neutron to the target without the creation and subsequent decay of a compound nucleus. This leads to proton angular distributions that are either forward or backward peaked in contrast with the expectations of the CN theory that predicts symmetric angular distributions.

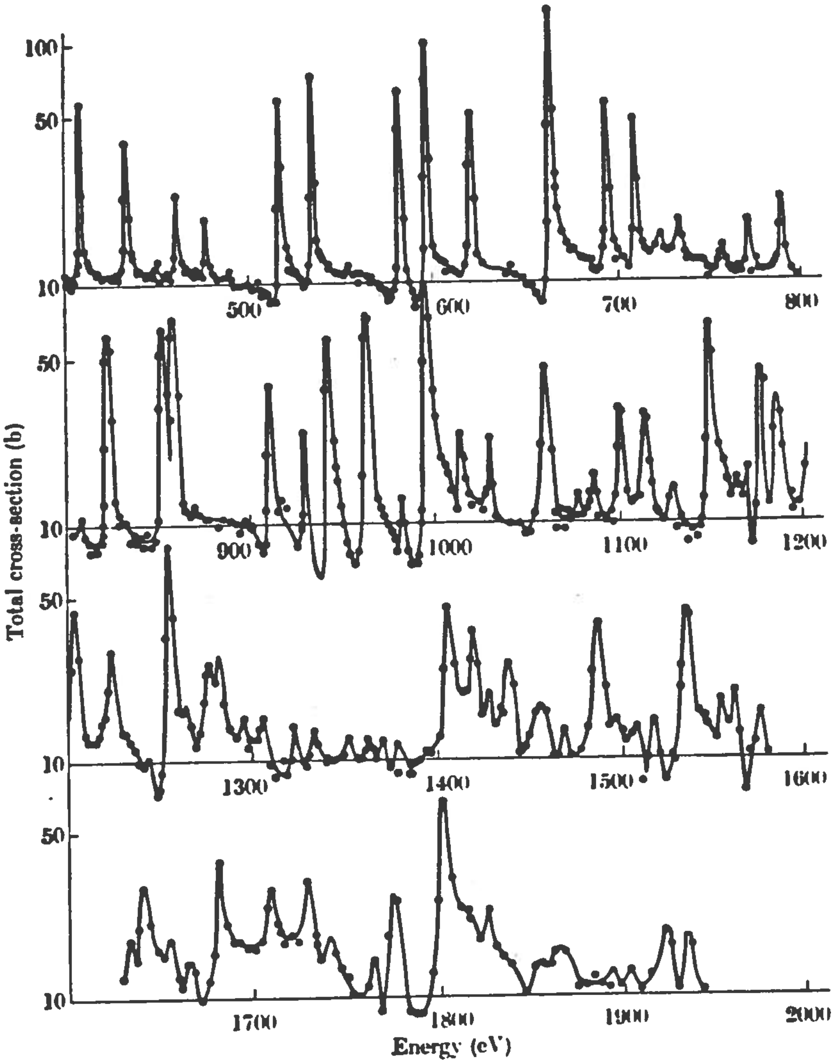
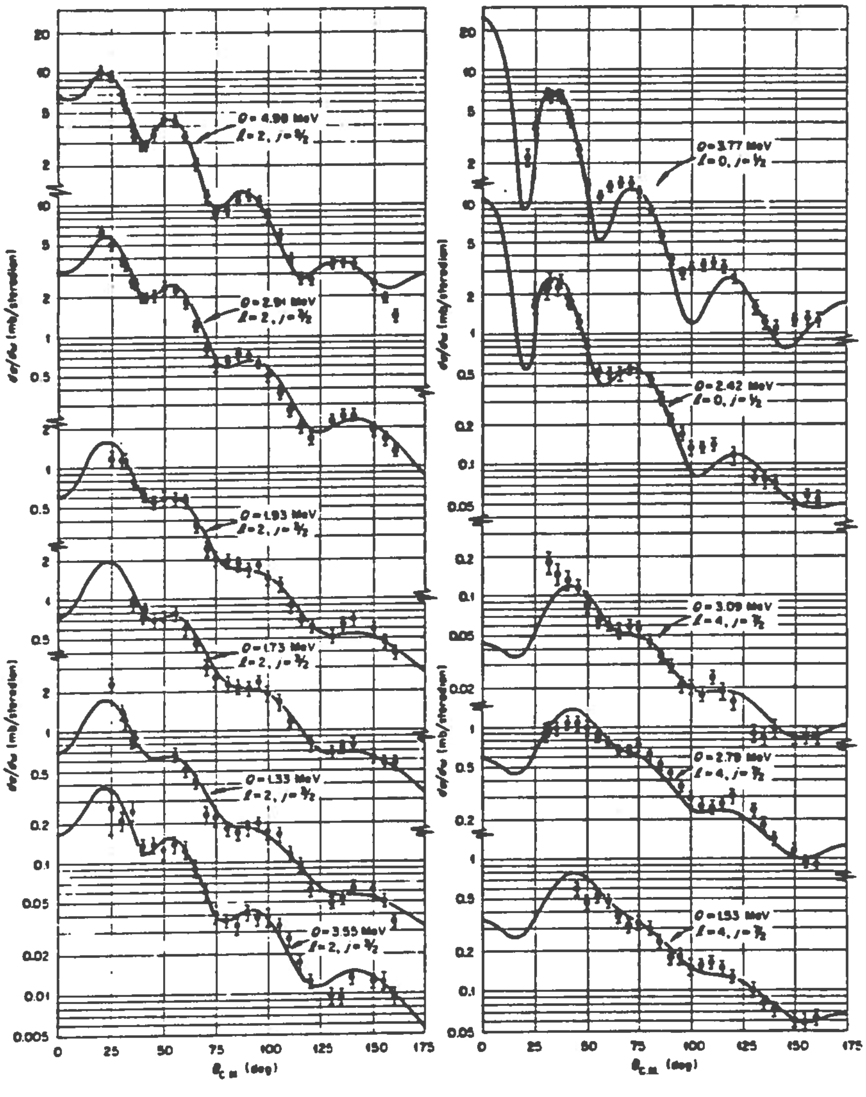

Total cross-section (b)
100
50
10
500
600
700
800
50
10
900
1000
1100
1200
50
10
1300
1400
1500
1600
50
10
1700
1800
1900
2000
Energy (eV)

Fig. 1

More generally the term *direct reactions* is now applied to all processes that directly connect the initial and final states in a nuclear reaction without formation of an intermediate compound system. Typical examples of direct processes are *elastic scattering* which can be explained by a very simple model (the optical model) in which the interaction between the incident nucleon and the nucleus is given by a one-body potential, *inelastic scattering* that predominantly excites collective states, *charge-exchange reactions* in which the emitted particle has the same mass as the projectile, so that the net effect is the transfer of charge. Other direct reactions are the *stripping* and $pick - up$ reactions that involve the transfer of one or more nucleons. Typical exemples of angular distributions of particles produced in direct reactions are shown in Fig. 2 where the measured and the calculated angular distributions of protons emitted in the $^{90}Zr(d,p)^{91}Zr$ reaction at 12 MeV incident energy are compared (Dickens *et al* 1967).

Direct reactions are a powerful spectroscopic tool that allows one to study the configuration of the nuclear levels that may be excited in a given reaction and to determine their spin and parity. For instance, a (d,p) reaction is only possible if the single-particle state capturing the neutron is wholly or partially empty, so the observed cross-section is a measure of its emptiness. Conversely, the cross section of the inverse (p,d) pick-up reaction enables the occupation probabilities of the neutron single particle states in the target nucleus to be determined.

CN formation and decay and direct processes give rise to a very different energy and angular dependence of the cross-sections, so they may be easily differentiated: (a) the process of evaporation of particles from the CN favours emission of low energy particles, that is neutrons of a few Mev and charged particles with energy exceeding by a few MeV that of the Coulomb barrier. On the contrary a direct process favours emission of high energy particles with the excitation of low energy states of the residual nucleus., (b) the long lifetime of the CN, implying loss of memory of the way it was formed, together with angular momentum conservation, leads to symmetric angular distributions while direct processes usually show forward-peaked angular distributions of a diffractive character., (c) CN reactions are not selective,



Q = 4.98 MeV
ℓ = 2, j = 5/2
Q = 2.94 MeV
ℓ = 2, j = 3/2
Q = 1.93 MeV
ℓ = 2, j = 3/2
Q = 1.73 MeV
ℓ = 2, j = 3/2
Q = 1.33 MeV
ℓ = 2, j = 3/2
Q = 3.55 MeV
ℓ = 2, j = 3/2
Q = 3.77 MeV
ℓ = 0, j = 1/2
Q = 2.42 MeV
ℓ = 0, j = 1/2
Q = 3.09 MeV
ℓ = 4, j = 7/2
Q = 2.79 MeV
ℓ = 4, j = 7/2
Q = 1.53 MeV
ℓ = 4, j = 7/2
dσ/dω (mb/steradian)
θ_C.M. (deg)
0
25
50
75
100
125
150
175
20
10
5
2
1
0.5
0.2
0.1
0.05
0.02
0.01
0.005

Fig. 2

populating states of the residual nucleus with intensity proportional to their statistical weight, while direct transitions connect states of the target and residual nucleus with a definite structural relationship and, as a consequence, are strongly selective., (d) finally the excitation functions of CN processes, in good energy resolution experiments, show a rapid variation with the energy; at low energy a resonance behaviour arising from the excitation of well separated CN states, and at higher energies a fluctuating behaviour due to the random character of the reduced widths contributing to the transition amplitude, as shown by Ericson (1963b). The excitation functions of direct processes show a weak variation with energy sometimes including broad resonances.

Starting from 1947, the accumulation of experimental results indicated the presence of processes different in character from either the direct or the CN mechanisms that are now known as pre-compound or pre-equilibrium reactions. A typical exemple of processes of this type is shown in Fig. 3 where the spectra of protons inelastically scattered by ^{54}Fe at various incident energies are reported (Bertrand and Peelle 1973). One may see the low energy evaporative peak typical of CN processes, the high energy resolved peaks for direct transitions to low excitation energy states, the intermediate structureless pre-compound component.

To explain these processes Serber (1947) made the hypothesis that the interaction of a high energy nucleon with the nucleus could, to a good approximation, be described in terms of two-body interactions of the projectile with nucleons of the target initiating a cascade of nucleon-nucleon interactions which spreads the projectile energy among an ever increasing number of nucleons. Thus one may distinguish two stages in the reaction: a *fast stage*, corresponding to the cascade of nucleon-nucleon interactions, during which either the incident particle or a particle struck near the nuclear surface may emerge from the nucleus and a ***slow stage*** when, at the end of the intranuclear cascade, a fully equilibrated nucleus is created which further de-excites by the usual evaporation process. In this work Serber first suggested that the nucleon-nucleon collisions inside the nucleus could be described as free collisions provided that the degeneracy of nuclear matter be taken into account

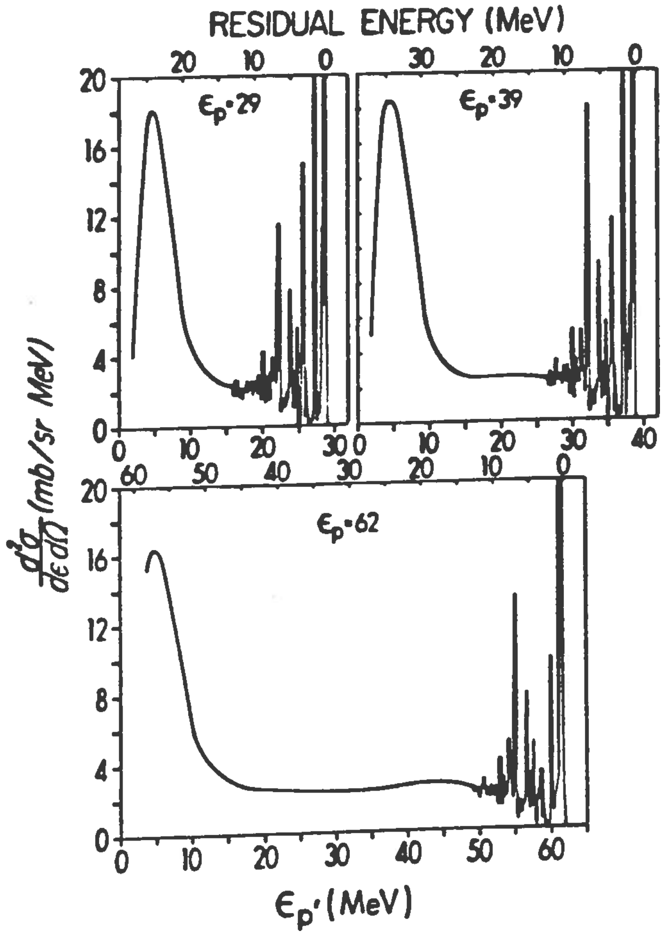

RESIDUAL ENERGY (MeV)
20
10
0
30
20
10
0
ϵ_p=29
ϵ_p=39
20
16
12
8
4
0
$\frac{d^2\sigma}{d\epsilon d\Omega}$ (mb/sr MeV)
0
10
20
30
0
10
20
30
40
60
50
40
30
20
10
0
ϵ_p=62
20
16
12
8
4
0
0
10
20
30
40
50
60
$\epsilon_{p'}$ (MeV)

Fig. 3

since, if nuclear matter is represented as a degenerate Fermi gas, collisions with small momentum transfer will be discouraged, since these tend to lead from an occupied state to another already occupied. The connection between C N, direct reactions and pre-equilibrium processes is illustrated in Fig. 4.

In subsequent years many calculations based on this reaction mechanism were discussed in the literature. Among these, one by Goldberger (1948) in which for the first time the average nucleon-nucleon cross-section in nuclear matter, $\overline{\sigma}$, was evaluated by averaging the free interaction cross-section over the momentum distribution of the struck nucleons, taking approximately into account the effect of Pauli principle. These papers had a great influence on the phenomenological pre-equilibrium models that will be discussed later and were also the basis for the semi-classical optical model in which the imaginary part of the optical potential is expressed in terms of $\overline{\sigma}$, the nucleon velocity v and the nuclear density ρ as

$$W(r) = \frac{\hbar}{2}v(r)\rho(r)\overline{\sigma}(r). \tag{1}$$

Although these calculations are essentially qualitative they reproduce satisfactorily the main features of the phenomenological imaginary optical potential as shown in Fig. 5 where the imaginary optical potential which allows one to reproduce the angular distributions of protons elastically scattered by ^{58}Ni at various energies (full line) is compared to that evaluated by using relation (1) (Greenlees *et al* 1968). Note that at the lowest energy the proton angular distribution is sensible only to the high R values of the phenomenological imaginary optical potential.

In 1958 Metropolis *et al* published the result of Monte Carlo calculations of cascades of nucleon-nucleon interactions initiated by protons and neutrons with energy varying from 82 to 365 MeV. In these calculations the classical trajectories of the excited nucleons inside the nucleus are evaluated. Random numbers are used to decide if and where an interaction can take place and also the direction and momentum of the target nucleon as well as the directions and momenta of the particles after collision. A nucleon can be emitted when it reaches the nuclear surface with sufficiently high energy being directed outward and the cascade terminates when the energy of

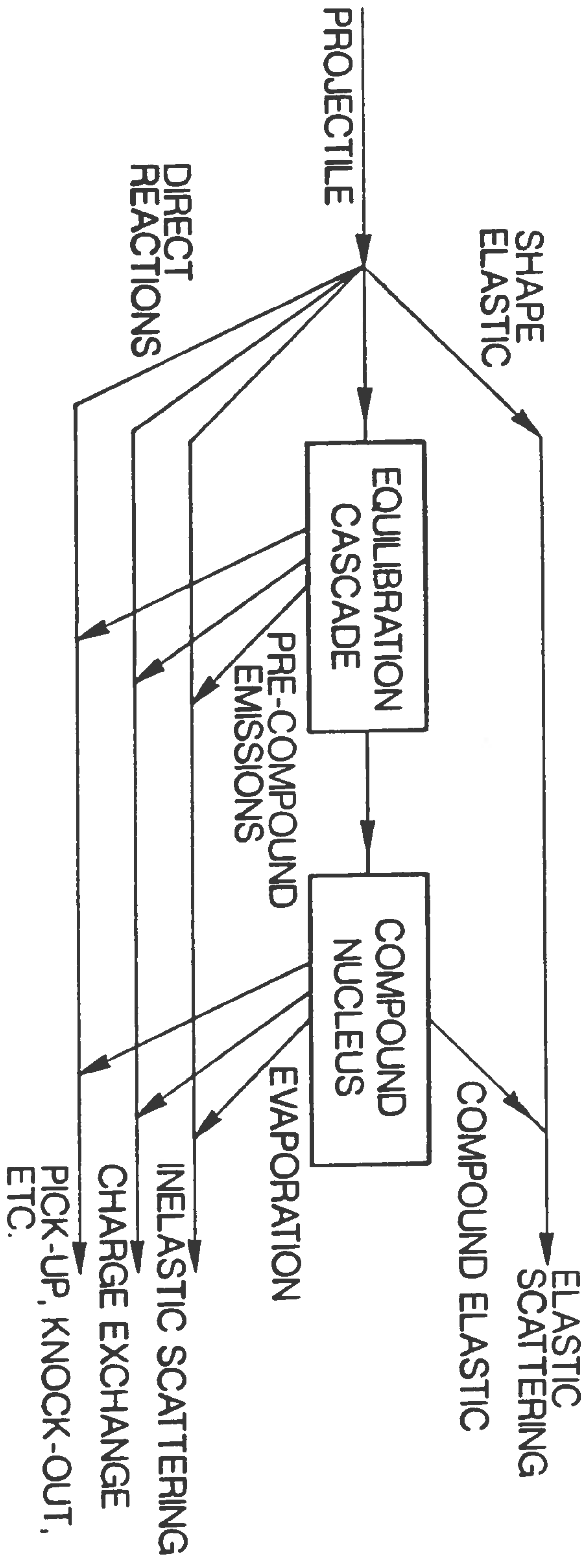

PROJECTILE
SHAPE ELASTIC
DIRECT REACTIONS
EQUILIBRATION CASCADE
PRE-COMPOUND EMISSIONS
COMPOUND NUCLEUS
EVAPORATION
COMPOUND ELASTIC
ELASTIC SCATTERING
INELASTIC SCATTERING
CHARGE EXCHANGE
PICK-UP, KNOCK-OUT, ETC.

Fig. 4

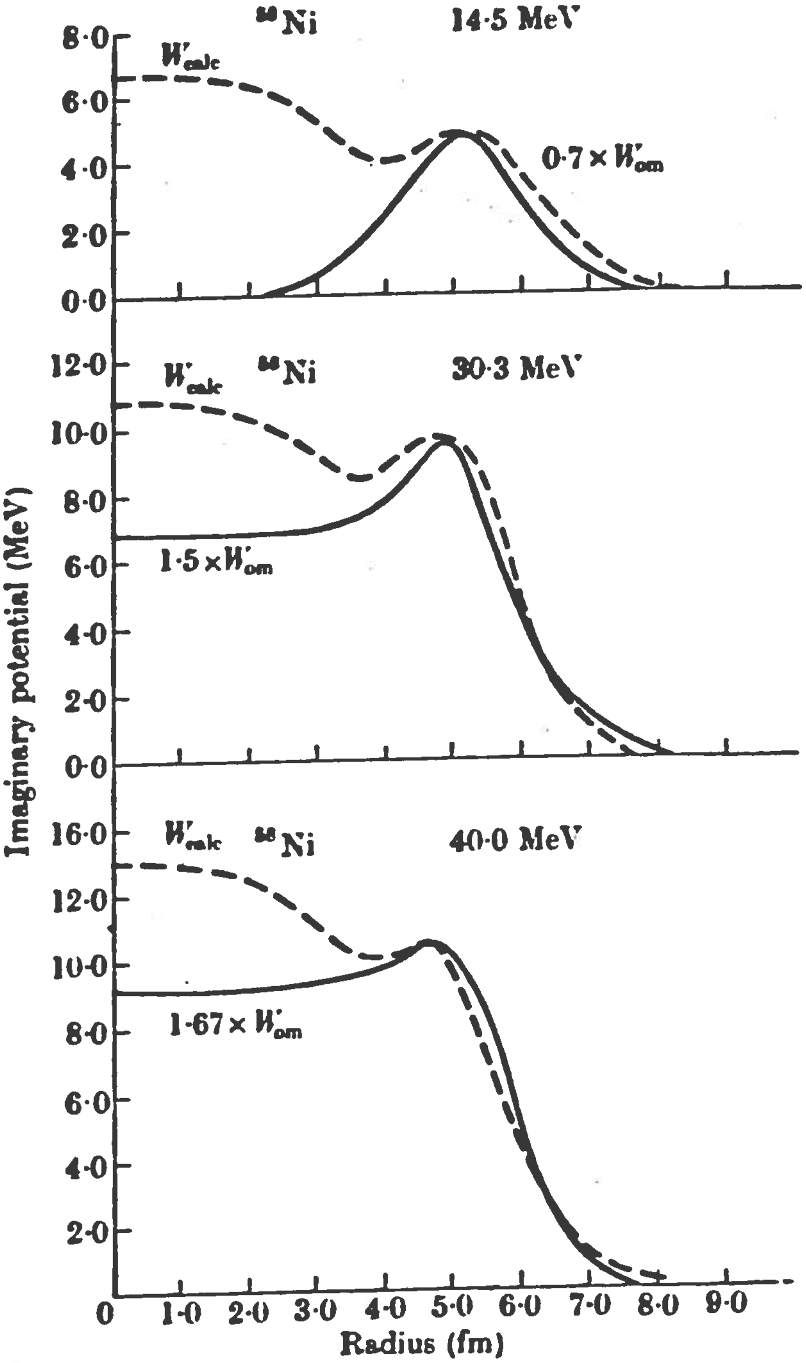

^{58}Ni
14·5 MeV
W'_{calc}
$0{\cdot}7 \times W'_{om}$
8·0
6·0
4·0
2·0
0·0
W'_{calc}
^{58}Ni
30·3 MeV
$1{\cdot}5 \times W'_{om}$
12·0
10·0
8·0
6·0
4·0
2·0
0·0
W'_{calc}
^{58}Ni
40·0 MeV
$1{\cdot}67 \times W'_{om}$
16·0
12·0
10·0
8·0
6·0
4·0
2·0
Imaginary potential (MeV)
0
1·0
2·0
3·0
4·0
5·0
6·0
7·0
8·0
9·0
Radius (fm)

Fig. 5

all nucleons falls below a predetermined value. Further de-excitation occurs through evaporations from the residual nucleus considered to be in statistical equilibrium. In these calculation the nucleus is described as a degenerate gas of nucleons. Free *p-p* and *n-p* cross-sections are used to predict the relative probabilities for each scattering event and when one or both the particles after a collision have energy lower than the Fermi energy the collision is forbidden. These calculations have been continuously improved through the years. One of the most sophisticated versions of them is the VEGAS code due to Chen *et al* (1968a,b, 1971). Typical events predicted by Monte Carlo calculations are shown in Fig. 6: (a) inelastic scattering following the interaction of the projectile with a nucleon near the surface of the target, (b) a cascade of nucleon-nucleon interactions with emission of pre-equilibrium particles, (c) formation of the compound nucleus. The Monte Carlo calculations, which reproduce the general features of the experimental data and are often in good quantitative agreement with them, even at rather low incident energies (Bertrand and Peelle 1973), are classical in nature (for instance, they totally neglect the uncertainty principle) and correspond to the kinetic theory limit of nuclear reaction theory.

Quantum mechanical theories of the progression toward the compound nucleus, started to appear in early sixties. In 1963 the concept of the *doorway state* was introduced by Block and Feshbach. The idea was that the excitation of the compound nucleus proceeds through states involving initially only a few degrees of freedom. They are intermediate in complexity between the single particle states described by the optical model and Bohr's compound nucleus states and their presence is revealed by resonances of width intermediate between that of the shape resonances, Γ_{SR}, typical of single particle states, and that of compound nucleus states, Γ_{CN}. They are observable in experiments made with an energy resolution ΔE satisfying the relation $\Gamma_{SR} \gg \Delta E \geq \Gamma_{CN}$ only if the doorway states are weakly coupled to the compound nuclear resonances. These doorway states were soon discovered in nucleon induced reactions (see, for instance, Elwyn *et al* 1965 and Singh *et al* 1965). In Fig. 7 the excitation functions for photocapture of protons by ^{27}Al with decay to the ground state of ^{28}Si are shown. With a very low energy resolution only a broad single particle

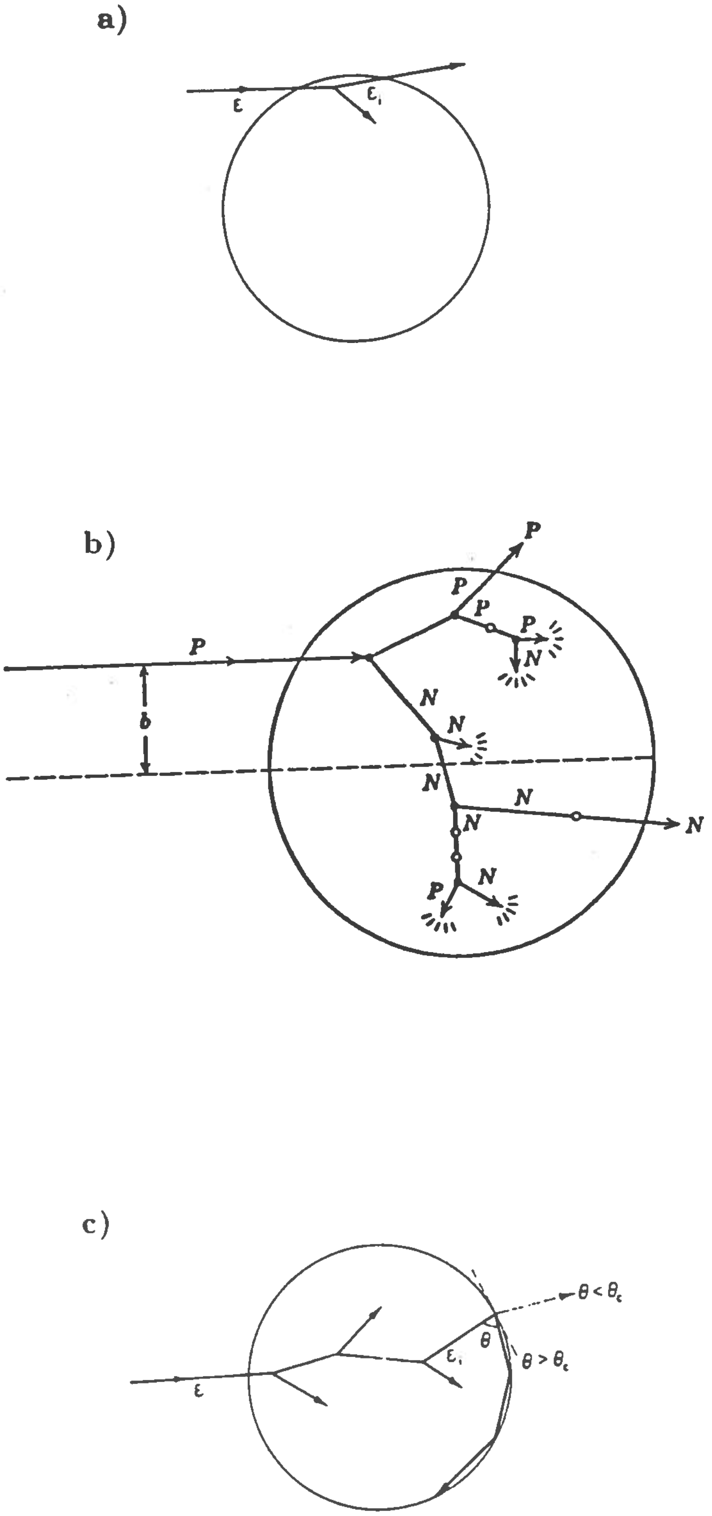


Fig. 6

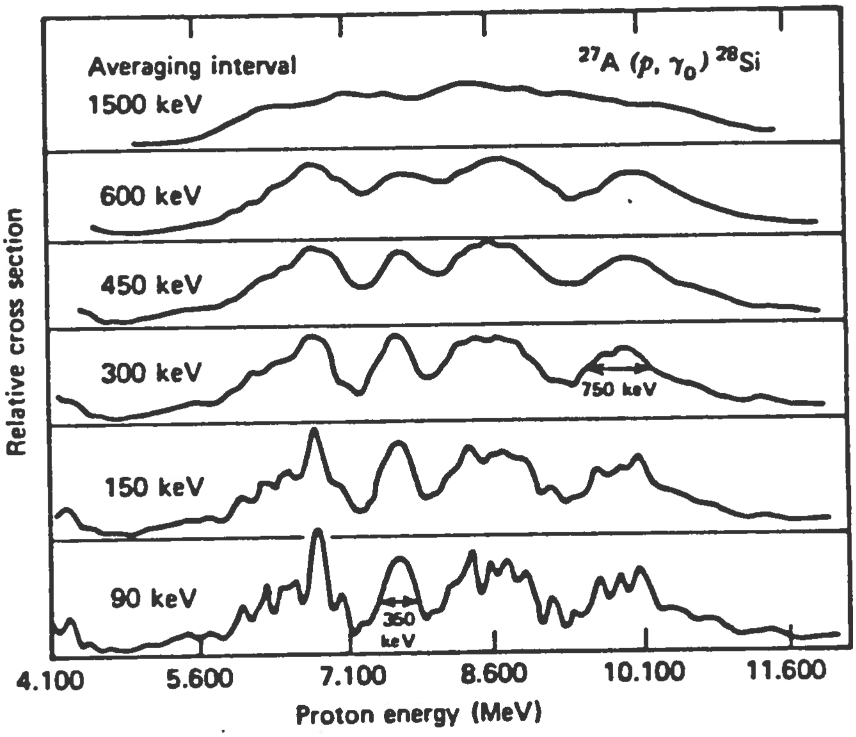

Averaging interval
1500 keV
^{27}A (p, γ_0) ^{28}Si
600 keV
450 keV
300 keV
750 keV
150 keV
90 keV
360 keV
Relative cross section
4.100
5.600
7.100
8.600
10.100
11.600
Proton energy (MeV)

Fig. 7

resonance is observed whose structure is revealed as the energy resolution improves. One may thus observe doorway state resonances of width $\Gamma_D \approx 750 KeV$ and the fine structure compound nucleus fluctuations.

The nature of the doorway states is not specified, they may be, for instance, particle-hole states or vibrational and rotational excitations. The former case is particularly significant for the influence it had on the subsequent development of precompound theories. If one assumes that the evolution from single particle to compound nucleus states is due to a residual interaction V_R which may be described by a two-body potential or a sum of two-body potentials, then in the case of a reaction induced by a nucleon these doorway states will have a two-particle-one-hole configuration. In 1967 Feshbach, Kerman and Lemmer "*emphasized that after averaging over the effect of complex states the doorway state resonances could be treated like ordinary compound nuclear resonances with the difference that the total width of the state contains not only the width due to the decay of the resonance into open channels but also an average "down" width describing the possibility of its decay into more complex states*"(Feshbach 1974). Even if the applicability of the theory to precompound emission was not indicated in these papers, this theory certainly had a great influence on further developments.

Another advance toward a more fundamental description of the de-excitation mechanism was made by Griffin (1966) who proposed a model "*for the formation and decay of the average compound-nuclear state in which a weak two-body residual interaction causes transitions among the eigenstates of an independent-particle Hamiltonian which lie in the region dE^* near the compound-nuclear excitation energy E^*. The independent-particle model states are classified according to the number of particles and holes (referred to indiscriminately as "excitons") excited from the even-even ground state. The limitation of a two-body interaction, that it can only effect energy-conserving transitions which change the number of excitons by 0 or ±2, is invoked and exploited to eliminate matrix elements which vanish identically....Decay is assumed to occur (in a very short time) to a state with outgoing particle of energy E_o and residual nucleus of energy U, whenever a nucleus makes a transition to* *an independent-particle state in which one exciton has energy E_o in the continuum, and the remaining excitons share the energy $U=E^*-(E_o+B)$, where B is the binding energy of the emitted particle...."*

Griffin also suggested that all the compound nucleus states corresponding to a given exciton number could be considered as equiprobable together with all the possible decay modes of this compound nucleus whose probability of occurrence, thus, depends on the available phase space. The Griffin compound nucleus is different from the Bohr compound nucleus since it is not in a statistical equilibrium and for this reason hereafter will be indicated as *composite* nucleus and its decay modes include both the emission of particles to open channels and the excitation of more complex configurations. A schematic representation of the first few stages of a reaction in the pre-equilibrium model is given in Fig. 8. The horizontal lines indicate equally spaced single particle states in the potential well and the particles are shown as solid circles.

The Griffin exciton model has been improved in the following years by several authors (Blann 1968, Williams 1970, Cline and Blann 1971, Braga-Marcazzan *et al* 1972, Birattari *et al* 1973, Gadioli, Gadioli Erba and Sona 1973 and Ribansky *et al* 1973, 1974). It was shown that it allows one to evaluate with a satisfactory accuracy the angle integrated cross sections of all the processes induced by an incident light particle with energy up to about 200 MeV using a unique set of basic parameters (necessary to evaluate the density of the n-exciton states of the composite and the residual nuclei, the level density of the equilibrated nuclei at the end of the de-exciting cascade, the nucleon-nucleon interactions inside the nuclear matter and so on). A typical exemple of the results obtained in the analysis of the data is shown in Fig. 9 where the measured and calculated excitation functions of the $^{48}Ti(p, xnyp)$ reactions are compared (Gadioli *et al* 1981).

A great effort to provide a formal justification of the assumptions made in the phenomenological theory of precompound reactions has been made by Wedenmüller and his collaborators. In a paper by Agassi, Weidenmüller and Mantzouranis (1975) a microscopic statistical model, based on the random-matrix theory of the nuclear Hamiltonian, was proposed. The exciton model hypotheses were shown to be ap-

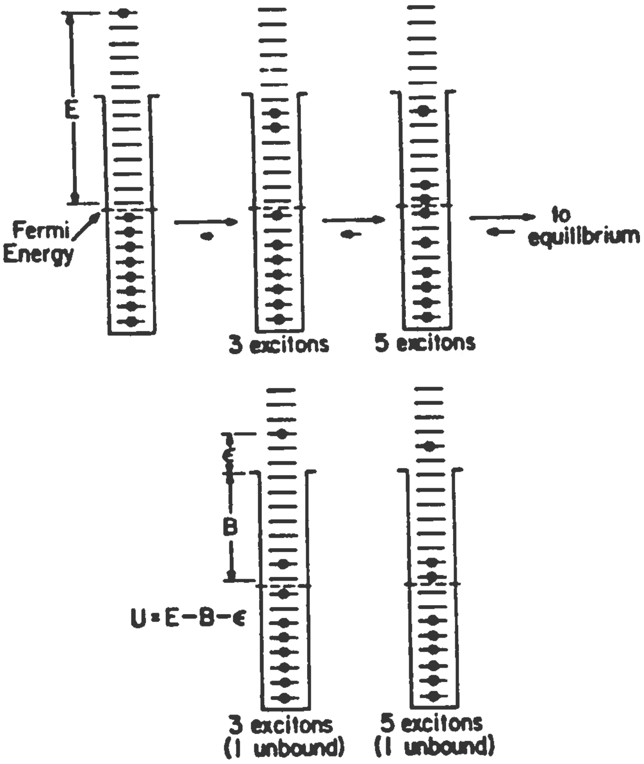
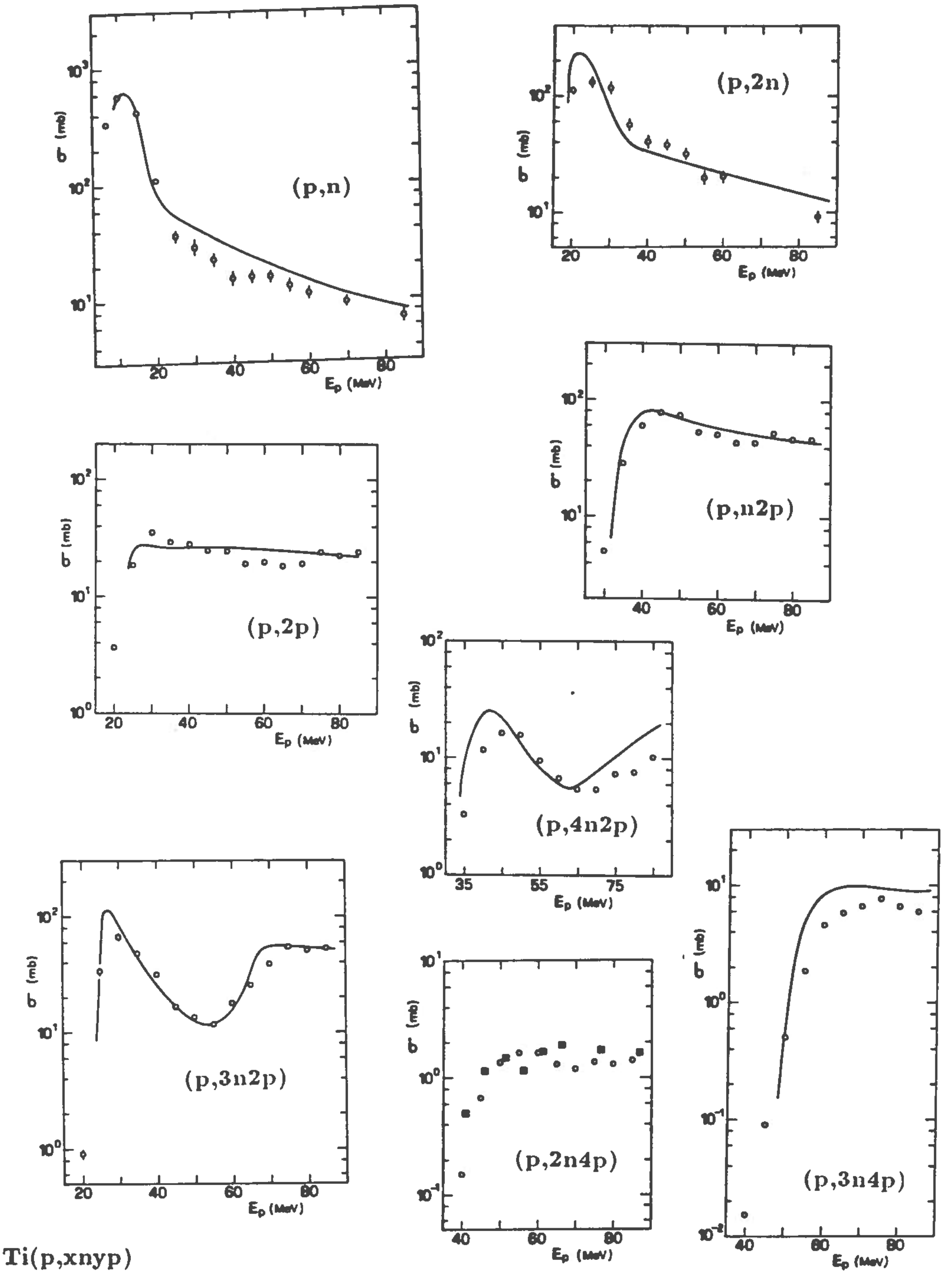

E
Fermi Energy
3 excitons
5 excitons
to equilibrium
ε
B
U=E−B−ε
3 excitons (1 unbound)
5 excitons (1 unbound)

Fig. 8



^{48}Ti(p,xnyp)

o, experiment

full line and black squares, theory

Fig. 9

proximately true by assuming that the mixing between the various classes of states reached during equilibration tends to zero, the weak coupling limit. Recently, using new mathematical techniques, Weidenmüller and coworkers (Verbaarschot *et al* 1985, Nishioka *et al* 1986, 1987) have extended the theory of precompound reactions beyond the weak-coupling limit.

A considerable effort has also been made, with a moderate success, for reproducing also the angular distributions of the particles emitted in a pre-equilibrium process. These are reproduced quite accurately up to ≈ 100°. Thereafter the calculations tend, generally, to underestimate the measured emitted particle yield. This seems to be due to quantum mechanical effects that it is not easy to incorporate within the model and prompted the development of the more formal theories of Feshbach, Kerman and Koonin (1980) (see also Feshbach 1973, 1974, 1977, 1979 and 1985) and Tamura and Udagawa (Feng *et al* 1976, Tamura and Udagawa 1977, Tamura *et al* 1977, 1981, 1982, Udagawa *et al* 1983).

Feshbach, Kerman and Koonin postulate two different statistical processes occurring in the course of the de-excitation of a nucleus before the evaporation stage, namely the statistical multistep compound (SMSC) and the statistical multistep direct (SMSD) processes. As in the case of the exciton model the reaction is considered to proceed through stages of increasing complexity. While in the case of SMSC, in each stage, all excited nucleons are bound, in SMSD, in each stage, at least one of the particles is in the continuum. The sequences of stages are called, respectively, the Q-chain in the case of SMSC and the P-chain in the case of SMSD. Two fundamental assumptions are made, the first, the *chaining hypothesis*, assumes that the residual interaction can induce transitions from the n-th stage only to the $(n \pm 1)$-th stages; the second is that the relative phases of certain matrix elements are assumed to be random. In the case of SMSC theory matrix elements involving different total angular momentum J, parity and the other quantum numbers required to specify a channel are assumed to have random relative phases, so that no interference terms remain upon averaging. The angular distributions generated by the multi-step compound process are therefore symmetric about 90° in the center of mass. By contrast, in the case of multi-step direct mechanism one assumes that the only matrix elements which interfere constructively upon averaging are those involving the same change in the momentum of the particle in the continuum. The *memory* of initial direction is therefore preserved and an asymmetric angular distribution results. The cross-section predicted by the multi-step direct reaction is added to that due to the multi-step compound reaction. It is given by a folding with respect to the momenta of the continuum particles of a number of direct transition probabilities equal to the number of stages. The final expression has a form consistent with a classical description of the reaction in momentum space and thus reduces to an approximation to a multiple scattering series at high energies. The multistep description of a nuclear reaction, according to FKK theory, is sketched in Fig. 10. Pre-equilibrium emissions can take place directly from each step of the P-chain, or indirectly from the Q-chain. The emissions from the Q-chain take place through states in the P-chain, and this can take place in three ways as shown in the figure. The more energetic particles come from the early stages of the chains and the less energetic from the later stages. This theory has been and is used with a considerable success in the analysis of the experimental data (Hodgson 1988).

The multistep direct reaction (MSDR) theory of Tamura and Udagawa, which constitutes a generalisation of the distorted wave Born approximation theory of direct reactions to deal with transitions to continuum states of the residual nucleus, is founded on the hypothesis that to reproduce the energy and angular distribution of the particles emitted in a multi-step process one needs only to evaluate the cross sections for transitions to pure shell model states and sum them incoherently. In fact, even if a multistep process feeds states of the residual nucleus that are a complicated superposition of pure shell model states, so in principle one should expect, in measured cross sections, interference between amplitudes for transitions to these different states, the assumption that the transition amplitudes to the shell model states are distributed statistically (with random signs and amplitudes) over the residual nucleus states, leads one to expecting that interference terms cancel to a large extent at high excitation energies where the residual nucleus states are highly overlapped.

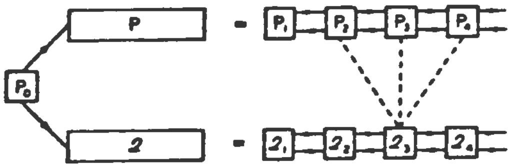

P
P₁
P₂
P₃
P₄
P₀
2
2₁
2₂
2₃
2₄

Fig. 10

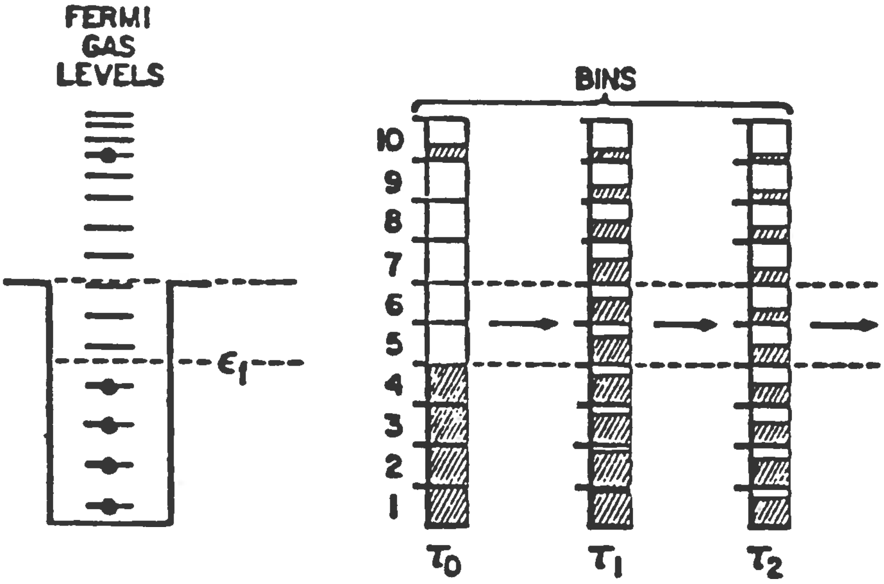

FERMI GAS LEVELS
ε₁
BINS
10
9
8
7
6
5
4
3
2
1
τ₀
τ₁
τ₂

Fig. 11

Assuming, further, that a given state can only be excited in a particular process, also interference between one-step and multi-step amplitudes need not to be considered and the contributions of the various steps to the total cross sections should be added incoherently. The total cross section for the transitions to the levels within a given excitation energy interval is proportional to their spectroscopic density that may be evaluated for any given process assuming a dominant reaction mechanism. For instance, in the case of (p,n) and (n,p) charge exchange reactions, the states that are expected to be excited are $1p-1h$ states and their spectroscopic density is expressed by means of the single particle response function. In case of inelastic scattering one may use either a spectroscopic density of this type or one evaluated with the random phase approximation (RPA) which describes rather accurately the low-lying collective states, including the giant resonance states. In the case of (p,α) reactions, assuming a dominant pick-up mechanism, the spectroscopic density is expressed by means of the overlap integral of three particle wave-functions and triton wave-functions of the correct spin-isospin nature and zero relative internal momentum.

Almost at the same time in which Griffin theory appeared, another major advance in the understanding of the mechanism of the de-excitation process was made by Harp, Miller and Berne (1968) (see also Harp and Miller 1971). In this theory (see Fig. 11, from Blann 1975) the nucleon states are classified according to their energy, ϵ, and divided into bins of width $\Delta\epsilon$. The number N_i of occupied states within each bin is equal to the product of the total number of states for that bin, g_i, times an occupation number $0\leq n_i\leq 1$. Nucleons in states within bins i and j may interact and scatter to states within bins l and m subject to the conservation of energy and the availability of unoccupied states in l and m . Unbound nucleons may also escape from the nucleus with energy $\epsilon_i'=\epsilon_i-\epsilon_F-B_i$ (ϵ_F and B_i are, respectively, the Fermi energy and the binding energy of the nucleon in the composite nucleus) thus contributing to precompound emission.

The relaxation of the nucleus, is described by a system of master equations. This approach has not been greatly used in the analysis of reactions induced by light ions where the exciton model represents an equivalent and much more versatile approach.

The opposite occurs in the case of the fusion and the quasi-fusion of two heavy ions where, once the two ions come in contact creating a di-nuclear system, the energy distribution between the excited particles and holes is very different from that expected by assuming the equiprobability of all the states that may be excited. This distribution is, in first approximation, that expected in the hypothesis of the coupling of the translational momentum of the nucleons of both the projectile and the target (due to their being part of a translating ion) and their Fermi motion momentum within each ion (Robel 1979, Bondorf *et al* 1980). Exemples of the expected occupation number distributions of the nucleon states of this intermediate system are given in Fig. 12 (from Fabrici *et al* 1989) where the neutron state distributions are shown in the case of $600MeV$ ^{20}Ne ions on ^{165}Ho (dashed line) and 800 MeV ^{40}Ar on ^{40}Ca (full line). Neutrons with energy in excess of, respectively, the dashed and the full arrows are in the continuum. The dotted line gives the occupation number distribution for a zero temperature Fermi gas. These distributions are characterised by holes, in addition to excited particles, whose number, which may be very large, and excitation energy depend on the projectile-target mass symmetry. The presence of these holes might give rise to a new effect, not present in light ion reactions, consisting in the excitation of particles to an energy higher than their initial energy, as the cascade of nucleon-nucleon collisions develops. Fabrici *et al* (1989b) have investigated quantitatively this effect. In case of very asymmetric systems like $C + Ho$ or $Ne + Ho$, the calculated spectrum of emitted particles closely resembles that obtained by means of a generalisation of the exciton model which consists in neglecting the hole excitations and in considering simply the unbound excited particles that in the case of very asymmetric systems are essentially the projectile nucleons. This demonstrates that for asymmetric systems like those above considered (asymmetry parameter $y = A_p/A_T \leq 0.13$, A_p and A_T are, respectively, the projectile and the target mass) the effect of the interactions between excited nucleons in the presence of deep holes is of negligible importance. Basically different is the result one finds when one considers a symmetric system ($y = 1$) such as ^{40}Ar + ^{40}Ca at $800MeV$. In this case, the maximum neutron energy at the beginning of the nucleon-nucleon interaction cascade, corre-

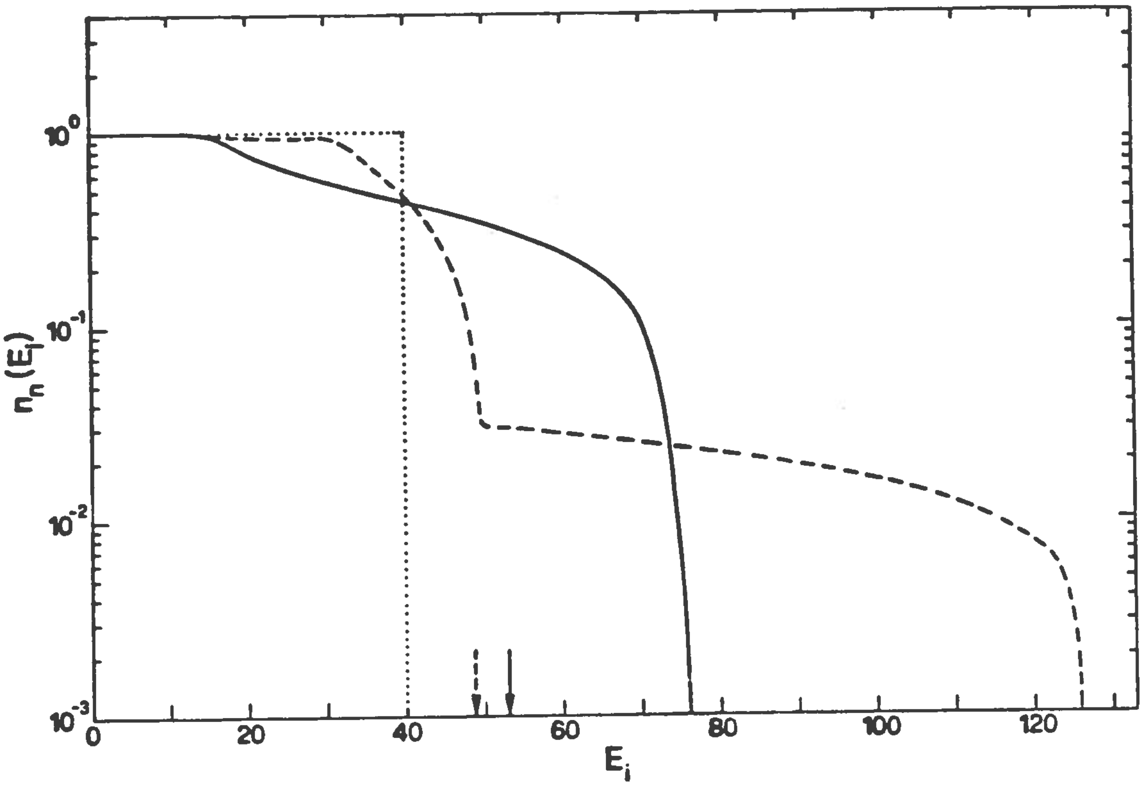


Fig. 12

sponding to the kinematical limit $E_K = (p_t + p_F)^2/2m_n - B - \epsilon_F$ (where p_t and p_F are the translational and the Fermi momentum; B, ϵ_F and m_n are, respectively, the neutron binding energy, the Fermi energy and the nucleon mass), is only $\approx$20.3 MeV, an energy considerably smaller than that to which the experimental neutron spectrum (measured by Rösch *et al* 1987) extends. Also considering the smearing due to the low energy resolution affecting these data, one cannot bring the calculations in agreement with the data in the absence of a considerable hardening of the spectrum due to nucleon-nucleon collisions. In fact, this hardening occurs as shown in Fig. 13 where the initial distribution of the neutrons with energy in the continuum (thick full line) is compared to that predicted at the end of the cascade of nucleon-nucleon interactions (thin full line). Fig. 14 shows that the spectra calculated using free nucleon-nucleon cross-sections (thin full line) and cross-sections scaled, respectively, by a factor 2 (dashed line) and 4 (dotted line) are in excellent agreement with the measured spectrum (thick full line).

In Fig. 15 are shown, for systems of different symmetry ((A) ^{32}S + ^{27}Al ($y = .84$, full line), (B) ^{32}S + ^{58}Ni ($y = .55$, dotted and dashed line), (C) ^{32}S + ^{120}Sn ($y = .27$, dotted line), (D) ^{32}S + ^{197}Au ($y = .16$, dashed line)), the calculated spectra of pre-equilibrium protons. The energy of the incident ^{32}S ion beam is in all cases equal to 679 MeV and, in the calculation, free nucleon-nucleon cross sections have been used. The arrows indicate, in each case, the highest energy of the protons at the beginning of the interaction cascade resulting from the kinematical coupling of the translational and internal momentum. The importance of the nucleon-nucleon interaction is shown by the fact that, in spite of the considerably different initial energy distributions the emitted particle spectra are very similar and, if there is a difference, the softest spectrum corresponds to the hardest initial distribution. In case (D) one may identify two components of the calculated spectrum, one corresponding to the particles emitted immediately, without any interaction, that ends at the energy corresponding to the arrow (D), and one of protons that acquired energy in at least one previous nucleon-nucleon interaction. The two contributions are barely observable in the case of ^{120}Sn, and cannot be separated in all the other cases.

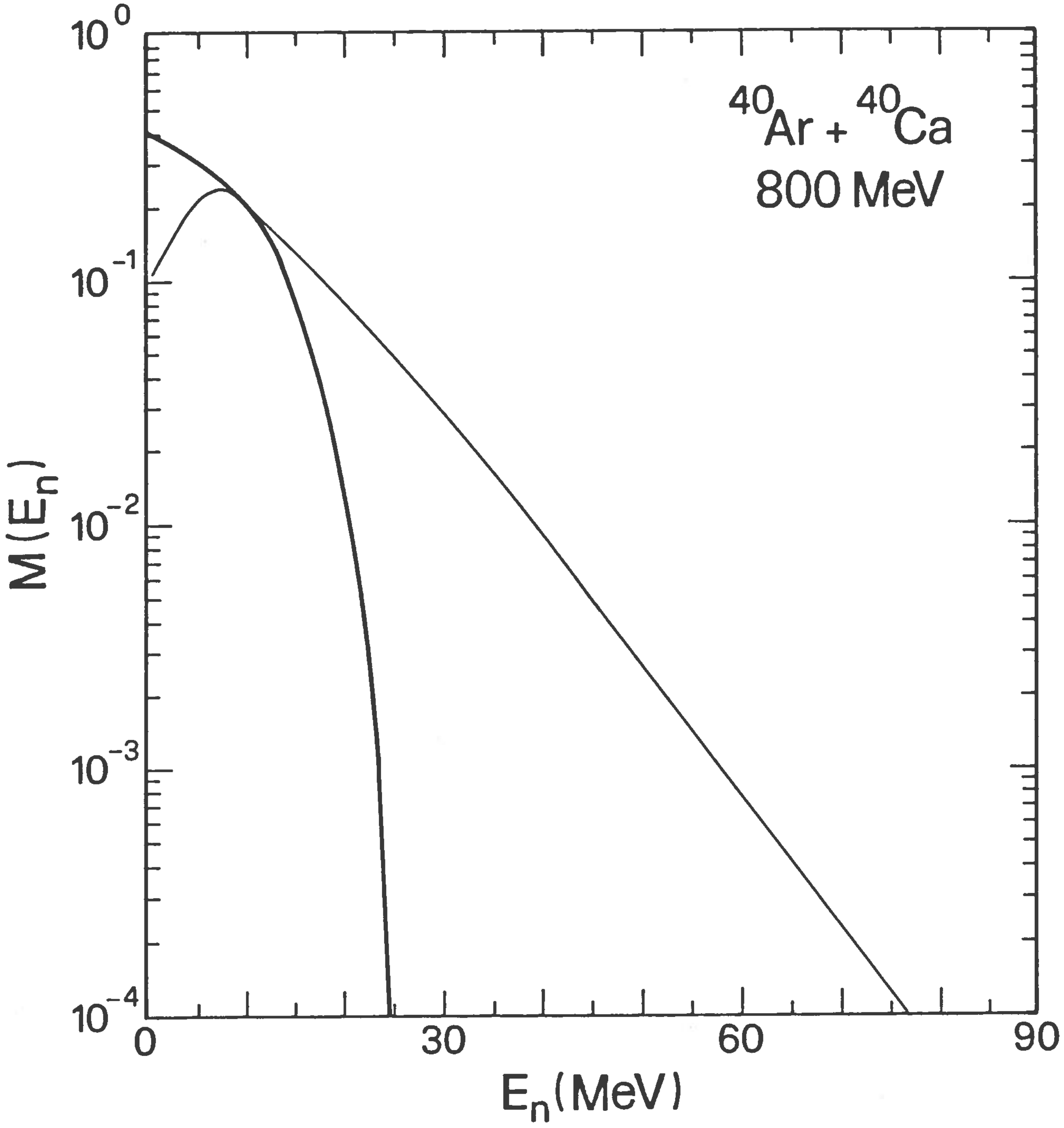


Fig. 13

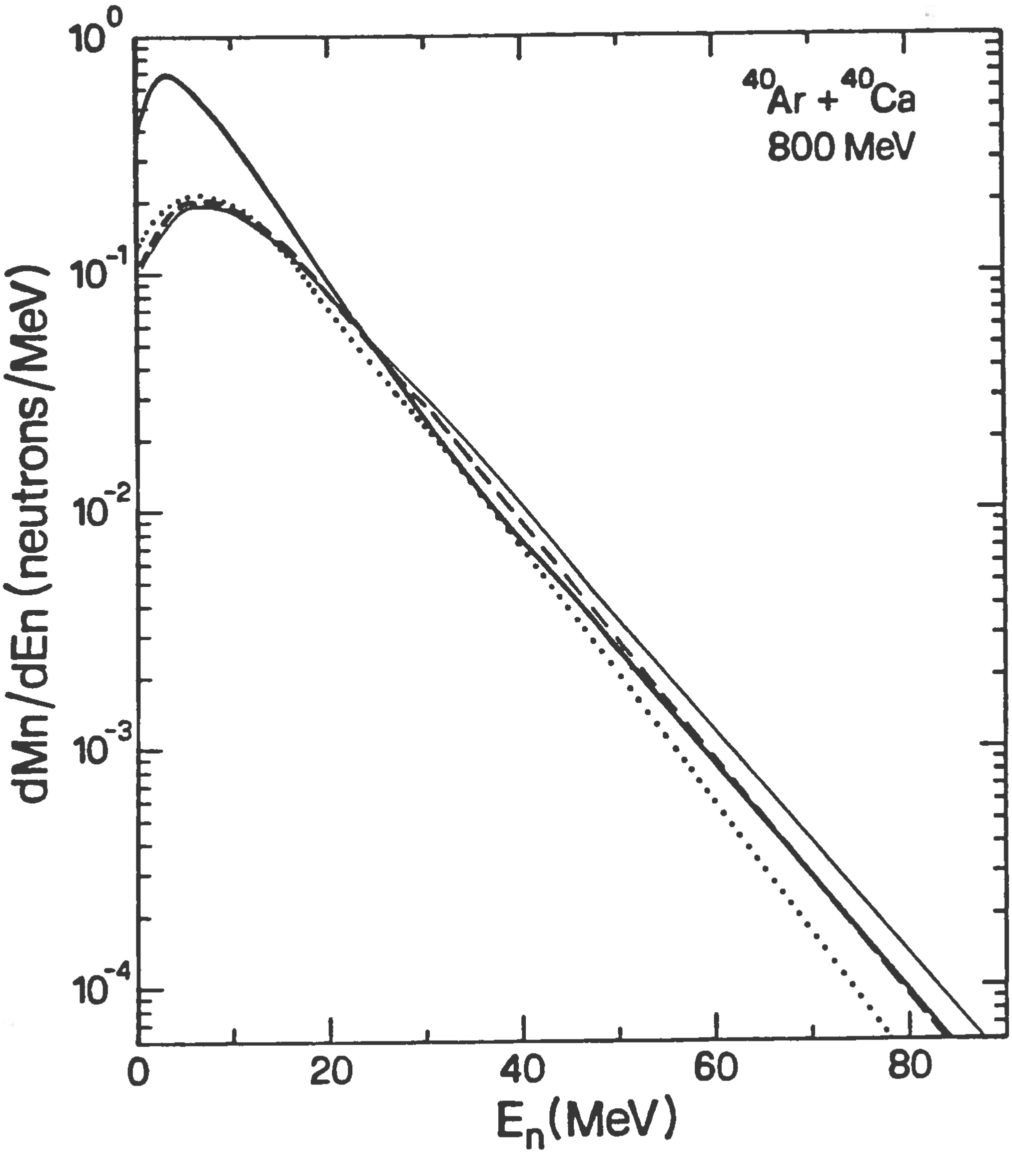
^{40}Ar + ^{40}Ca
800 MeV
dMn/dEn (neutrons/MeV)
10^0
10^{-1}
10^{-2}
10^{-3}
10^{-4}
0
20
40
60
80
E_n(MeV)

Fig. 14

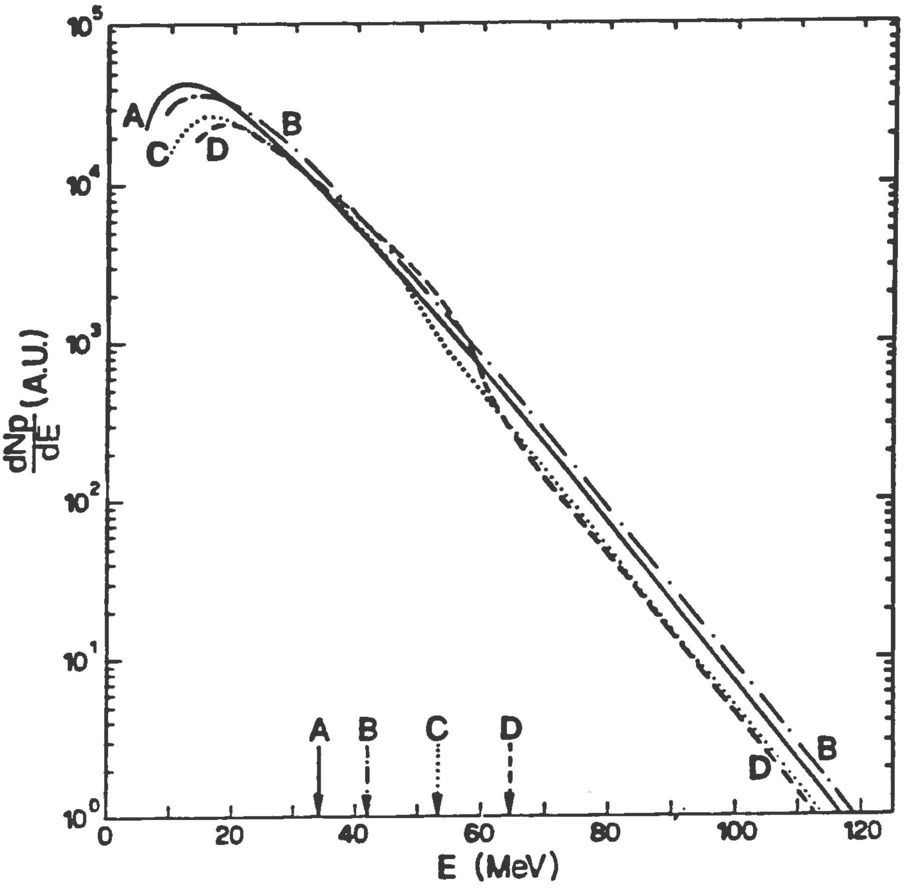

10^5
10^4
10^3
10^2
10^1
10^0
$\frac{dN_p}{dE}$ (A.U.)
A
B
C
D
A
B
C
D
D
B
0
20
40
60
80
100
120
E (MeV)

Fig. 15

The hardening of the nucleon spectrum obviously depends on the probability of interaction of two particles of sufficient high initial energies ϵ_i and ϵ_j which is proportional to the product of the corresponding number of occupied states $n_i g_i \cdot n_j g_j$ and the availability of deep holes at energy ϵ_l where one of the two particles may scatter, which is proportional to $(1-n_l)$. The comparison, in Fig. 12, of the initial occupation number distributions corresponding to asymmetric system $Ne + Ho$ and to the symmetric system $Ar + Ca$ shows that the probability of a nucleon-nucleon interaction producing nucleons with energy greater than that of the interacting nucleons may be orders of magnitudes greater in the symmetric case.

To conclude, starting from incident energies of the order of a few ten MeV/nucleon, nucleon-nucleon collisions play a major role in the equilibration process of the intermediate composite system created in the interaction of the projectile and the target both in the case of reactions induced by light projectiles and in heavy ion interactions. Phenomenological models based on simple statistical hypotheses made it possible to predict with a satisfactory accuracy the cross-section of the processes that occur during the development of the de-excitation cascade which follows the initial projectile-target interaction. More formal theories have also been proposed that aim to give a full quantum-mechanical description of these phenomena. All this work was necessary to understand how the compound nucleus state, in statistical equilibrium, predicted by Bohr in 1936 is reached and has greatly contributed to our present knowledge of the nuclear reaction mechanism.

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References

Agassi D, Weidenmüller H A and Mantzouranis G 1975 Phys. Rep. **22C** 147

Bertrand F E and Peelle R W 1973 Phys. Rev. **C8** 1045

Birattari C, Gadioli E, Gadioli Erba E, Grassi Strini A M, Strini G and Tagliaferri G 1973 Nucl. Phys. **A201** 579

Blann M 1968 Phys. Rev. Lett. **21** 1357

—— 1975 Ann. Rev. Nucl. Sci. **25** 123

Block B and Feshbach H 1963 Ann. Phys. (N. Y.) **23** 47

Bohr N 1936 Nature **137**, 344

Bondorf J B, De J N, Fai G, Karvinen A O T, Jakobsson B and Randrup J 1980 Nucl. Phys. **A333** 285

Braga Marcazzan M. G. *et al* 1972 Phys. Rev. **C6** 1398

Breit G and Wigner E 1936 Phys. Rev. **49**, 519

Brink D M and Stephen R O 1963 Phys. Lett. **5** 77

—— 1964a Phys. Lett. **8** 324

Brink D M *et al* 1964b Nucl. Phys. **54** 577

Chen K *et al* 1968a Phys. Rev. **166** 949

—— 1968b Phys. Rev. **176** 1208

—— 1971 Phys. Rev. **C4** 2234

Cline C K and Blann M 1971 Nucl. Phys. **A172** 225

Dickens J K, Perey F G and Silva R G 1967 *Int. Nucl. Phys. Conf.*, R L Becker, C D Goodman, P H Stelson and A Zucker Edts, Academic Press (N. Y.)

Elwyn A J *et al* 1965 Argonne Natl. Lab. Summary Report ANL-7081

Ericson T 1960 Phys. Rev. Lett. **5** 430

—— 1963a Phys. Lett. **4** 258

—— 1963b Ann. Phys. (NY) **23** 390

Fabrici E, Gadioli E, Gadioli Erba E, Galmarini M, Fabbri F and Reffo G 1989 Phys. Rev. **C40** XXX

Feng D H, Udagawa T and Tamura T 1976 Nucl. Phys. **A274** 262

Fermi E *et al* 1934 Proc. Roy. Soc. **A146**, 483

—— 1935a Proc. Roy. Soc. **A149**, 522

—— 1935b La Ricerca Scientifica **A6**, 544

Feshbach H, Kerman A K and Lemmer R H 1967 Ann. Phys. (N. Y.) **41** 230

Feshbach H 1973 *Proc. Int. Conf. Nucl. Phys.*, Munich, eds. J J DeBoer and H J Mang, vol. II, pg 631, North Holland Publ. Co., Amsterdam

—— 1974 Revs. Mod. Phys. **46** 1

Feshbach H 1977, 1979 *Proceedings of the 1st (2nd) Int. Conf. on Nucl. Reaction Mechanisms*, Varenna, Clued Publ. Co., Milano

Feshbach H, Kerman A K and Koonin S 1980 Ann. Phys. (N. Y.) **125** 429

Feshbach H 1985 Ann. Phys. (N. Y.) **159** 150

Firk F W K, Lynn J E and Moxon M C 1960 *Int. Conf. on Nucl. Structure,* Kingston, University of Toronto

Frisch R *et al* 1936a Nature **137**, 149

—— 1936b Nature **137**, 357

Gadioli E, Gadioli Erba E, Hogan J J, Burns K I 1981 Z. Phys. **A301** 289

Gadioli E, Gadioli Erba E and Sona P G 1973 Nucl. Phys. **A217** 589

Goldberger M L 1948 Phys. Rev. **74** 1268

Greenlees G W, Pyle G J and Tang Y C 1968 Phys. Rev. **171** 1115

Griffin J J 1966 Phys. Rev. Lett. **17** 478

Harp G D, Miller J M and Berne B. J 1968 Phys. Rev. **165** 1166

Harp G D and Miller J M 1971 Phys. Rev. **C3** 1847

Hauser W and Feshbach H 1952 Phys. Rev **87** 366

Hodgson P E 1988 *Proceedings of the Int. Conf. on Nuclear data for Science and Technology*, Mito (Japan), S. Igarasi Edt., Saikon Publ. Co. Ldt.

Metropolis N *et al* 1958 Phys. Rev **110** 185

Nishioka H, Verbaarschot I J M and Weidenmüller H A 1986 Ann. Phys. (N. Y.) **172** 67

Nishioka H, Weidenmüller H A and Yoshida S 1987 Ann. Phys. (N. Y.) **183** 166

Oppenheimer J R and Phillips M 1935 Phys. Rev. **48**, 500

Ribansky I *et al* 1973 Nucl. Phys. **A205** 545

—— 1974 Nucl. Phys. **A226** 347

Robel M C 1979 Ph. D. Thesis, LBL-8181, *unpublished*

Rösch W, Richter A, Schrieder G, Gentner R, Keller K, Lassen L, Lücking W, Schreck R, Cassing and W, Gemmeke H 1987 Phys. Lett. **197B** 19

Serber R 1947 Phys. Rev. **72**, 1114

Singh P P *et al* 1965 Nucl. Phys. **65** 577

Szilard L 1935 Nature **136**, 849

Tamura T and Udagawa T 1977 Phys. Lett. **71B** 273

Tamura T, Udagawa T, Feng D H and Kan K K 1977 Phys. Lett. **66B** 109

Tamura T, Lenske H and Udagawa T 1981 Phys. Rev. **C23** 2769

Tamura T, Udagawa T and Lenske H 1982 Phys. Rev. **C26** 379

Udagawa T, Low K S and Tamura T 1983 Phys. Rev. **C28** 1033

Verbaarschot I J M, Weidenmüller H A and Zirnbauer M R 1985 Phys. Rep. **129** 367

Weisskopf V F and Ewing D H 1940 Phys. Rev. **57** 472

Williams F C jr. 1970 Phys. Lett. **31B** 184