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Critical Review of Precompound Models with Consideration of the Contribution to the "International Nuclear Model and Code Comparison on Pre-Equilibrium Effects", organized by the NEA Data Bank 1983/84

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Abstract

It is shown that the desired predictive capability of most of the commonly used precompound formalisms to calculate nuclear reaction cross-sections is seriously reduced by too much arbitrariness of the choice of parameters. The origin of this arbitrariness is analyzed in detail and improvements and alternatives are discussed.

Kritische Besprechung der Precompound-Modelle unter Berücksichtigung des Beitrags zum "internationalen Kern-Modell- und Code-Vergleich von Vor-Gleichgewichtseffekten", veranstaltet von der NEA Data Bank 1983/84

Es wird hier gezeigt, daß die erwünschten Vorhersagensmöglichkeiten der meisten oft gebrauchten Precompound-Formalismen zur Berechnung von Kernreaktionsquerschnitten wegen zuviel Willkür der Parameterwahl erheblich eingeschränkt sind. Der Ursprung dieser Willkür wird analysiert und Verbesserungen und Alternativen werden diskutiert.

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Introduction

In recent years many measured data for secondary energy and angular dependent nuclear reaction cross-sections could be understood as representing events which occur during the equilibration process on the way until the compound nuclear states are rechead. The formal developments presenting this understanding seemed also to provide the necessary tools to calculate the considered cross-sections. But apparently it is overlooked quite often that there are important quantities occuring in most of the considered formalisms which have to be treated as parameters because they are too difficult to calculate and what is obtained is more a fit rather than a genuine predictive calculation. It is the purpose of this paper to show this in detail in order to obtain a help for a step on the way towards a more complete theory.

Sketch of the formalism

It is usually assumed that the nuclear reaction cross-sections split into a pre-equilibrium and an equilibrium component according to

$$(1) \qquad \frac{d\sigma(\varepsilon_i,\varepsilon_j)}{d\varepsilon_j} = \left(\frac{d\sigma(\varepsilon_i,\varepsilon_j)}{d\varepsilon_j}\right)_{preq.} + \left(\frac{d\sigma(\varepsilon_i,\varepsilon_j)}{d\varepsilon_j}\right)_{eq.}$$

where ε_i is the energy of the incident and ε_j the energy of the emitted particle. This additive splitting according to eq.(1) uses to be verified by means of the solutions of the following set of the so called master equations

$$(2) \qquad \frac{dP(n,t)}{dt} = P(n-2n,t)\lambda_+^{n-2,n} + P(n+2,t)\lambda^{n+2,n} - P(n,t)\,(\lambda_+^{n,n+2} + \lambda_-^{n,n-2} + L(n,E))$$

describing the evolution in time t of the probability P(n,t) that a certain total number n of particles and holes

$$(3) \qquad n = p + h$$

of the nuclear Fermi sea is excited. Cline and Blann /1/ have constructed

this set of master equations as a set of genuine balance equations describing the balance between the gains and the losses of probability for excitation of the n so called excitons. These gains and losses are caused by transition probabilities per unit time $\lambda_{\pm}^{n,n\pm2}(E)$ for creation or destruction of one particle-hole pair and by the total emission probability per unit time L(m,E) of a particle from a n-exciton state. Both $\lambda_{+}(E)$ and L(n,E) depend on the excitation energy E.

If we now consider t=0 as the time at which the reaction has started then the time T(n,E) spent by the composite nucleus in the n-excition state obviously is

$$T(n,E) = \int_0^\infty P(n,t)dt \tag{4}$$

Moreover we write as $W_j(n,E,\varepsilon_j)$ the probability per unit time for a particle of type j to be emitted with energy ε_j from an n-exciton state of excitation energy E. Thus

$$L(n,E) = \sum_j \int_0^{E-B_j} W_j(n,E,\varepsilon_j)d\varepsilon_j \tag{5}$$

where B_j is the binding energy of the particle of type j.
With the quantities T(n,E) and $W_j(n,E,\varepsilon_j)$ of (4) and (5) we obtain as the total cross section for emission of a particle of type j with energy between ε_j and $d\varepsilon_j$ by an impact of a particle of type i of energy ε_i

$$\frac{d\sigma_{ij}(\varepsilon_i,\varepsilon_j)}{d\varepsilon_j} = \sigma_{ci}(\varepsilon_i)\sum_{n=n_o} W_j(n,E,\varepsilon_j)T(n,E). \tag{6}$$

$\sigma_{ci}(\varepsilon_i)$ in (6) is the cross section for the formation of the composite system by the incoming particle i of energy ε_i. The summation is taken over all exciton states until the equilibrium is reached where n_o is the initial exciton number corresponding to the initial condition

$$P(n,o) = \delta_{nn_o} \tag{7}$$

According to Cline and Blann /1/, Ribanský, Oblozinský and Běták/2/, Wu and Chang /3/ and Dobeš and Běták /4/ the solution of (2) can very well be

approximated by an analytic closed-form expression. A corresponding approximate closed-form expression can consequently also be obtained for T(n,E) of eq.(4) which according to Dobeš and Běták /4/ can be written:

$$T(n,E) = T^{o}(n,E) + \alpha\omega_n \tag{8}$$

where

$$T^{o}(n,E) = \tau(n,E)\left[\prod_{\substack{i=n_o \\ \Delta n=2}}^{n-2} \lambda_{+}^{n,n+2}(E)\,\tau(i,E) + \delta_{nn_o}\right] \tag{8a}$$

with

$$\tau(n,E) = \left[\lambda_{+}^{n,n+2}(E) + \lambda_{-}^{n,n-2}(E) + L(n,E)\right]^{-1} \tag{8b}$$

and

$$\alpha = \frac{(1-\sum_n T^{o}(n,E)L(n,E))}{\sum_n L(n,E)\omega_n} \tag{8c}$$

The quantities $\lambda_{\pm}^{n,n\pm2}$ and L(n,E) in (8)-(8c) are those of (2) and (5). $T^{o}(n,E)$ in (8) is that part of the time integral (4) which goes from 0 until the equilibrium distribution of the exciton states at the equilibrium time t_{eq} is reached. For $t > t_{eq}$ this equilibrium distribution of the exciton states of course does not change anymore. It has to be taken proportional to the exciton state densities ω_n according to the postulate of equal a priori probability as has been pointed out by Cline and Blann /1/, Ribanský, Obložinský and Běták /2/ and Dobeš and Běták /4/. This has been used in the second term of (8).

In all of the recent work ω_n is expressed by the Ericson formula corrected by Williams /10/ to account for the Pauli principle. Thus

$$\omega_n = g\,\frac{(gE - A_{p,n})^{n-1}}{p!h!(n-1)!} \tag{9}$$

with the correction term $A_{p,n}$ due to the Pauli principle. $A_{p,n}$ was correctly presented by Williams /10/ only for the case p=h. To get an entirely correct

expression for $A_{p,h}$ several papers have been published (/5/,/6/,/7/ and /8/).
But none of them has presented the correct expression also for $p \neq h$.
The Williams-expression is /5/:

$$A_{p,h} = \frac{1}{4}(p^2 + h^2 + p - 3h) \tag{9a}$$

The correct expression has recently been found by Anzaldo /9/ as:

$$A_{p,h} = ph - \frac{1}{4}\left[p(p+1) + h(h+1) \right] \tag{9b}$$

On the other hand it could be shown /6/ that neglect of $A_{p,h}$ would not matter very much especially for the case where only one nucleon is incident or emitted and provided the excitation energy is not too small according to

$$E \gg \frac{A_{p,n}}{g} \tag{10}$$

But for more than one incident or emitted nucleon such as also for the case of α-particles or heavy ions the contribution of $A_{p,n}$ could become important. Note that (9) is based on the constant single-particle level density g taken at Fermi energy.

The exciton state density ω_n of (9) is of course also a factor in the expression $W_j(n,E,\varepsilon_j)$ of (5) and (6) for the particle emission probability per unit time. From the principle of detailed balance Cline and Blann /1/, /10/ have obtained the expression

$$W_j(n,E,\varepsilon_j) = \frac{2s_j+1}{\pi^2\hbar^3}\ \mu_j\varepsilon_j\sigma_{ci}Q_j(p)\ \frac{\omega_{n-1}(U)}{\omega_n(E)} \tag{11}$$

where s_j and μ_j are spin and mass of the emitted particle, U the excitation energy of the residual nucleus which for nucleons incident and emitted is

$$U = \varepsilon_i - \varepsilon_j \tag{12}$$

and $Q_j(p)$ is a combinatorial factor by which the proton-neutron distinguishability and more general the emitted particle type weighting is taken into

account to make it possible to use the one-Fermion type density of exciton states. By inserting (8), (9) and (11) into eq.(6) we obtain the additive splitting of eq.(1) where in the equilibrium term the denominator of (11) is cancelled by the ω_n of the second term of (8) and the remaining level density factor of the equilibrium term becomes

$$\omega(U) = \sum_n \omega_{n-1}(U) = \frac{\exp\{2\sqrt{\pi^2 gU/6}\}}{\sqrt{48U}} \tag{13}$$

The last expression of (13) has been obtained by Williams /5/ showing that the contributions from the Pauli principle correction term $A_{p,h}$ in (9) cancel in the summation of (13). Thus with (13) the one-Fermion type level density expression of Bethe for the free Fermi gas has been obtained in the equilibrium term of (1).

Transition rate problem

The most crucial quantities of the above sketched formalism are the transition rates $\lambda_{\pm}^{n,n+2}$ introduced with the master equation (2). After the first rough estimates of Griffin /1a/ and Blann /1b/ the following Golden Rule expression was stated by Williams /11/

$$\lambda_{\pm}^{n,n+2} = \frac{2\pi}{\hbar} \overline{|M|^2}\, \omega_{\pm}^{n,n\pm2} \tag{14}$$

where the square of the matrix element M is averaged over the indicated transitions. Correspondingly $\omega_{\pm}^{n,n\pm2}$ are exciton state densities taken for these transitions. The $\omega_{\pm}^{n,n\pm2}$ have first been calculated by Williams /11/ from the Ericson formula without and by Cline /6/ and by Obložinský, Ribanský and Běták /12/ from the Ericson formula with the Pauli correction term of Williams /5/ (see eq. (9) and the following text).
In addition the proton-neutron distinguishability has been taken into account by the above mentioned authors /12/. It amounts to a factor $\frac{1}{2}$ with which the expressions of the previous authors /11/, /6/ have to be multiplied.
The expressions thus obtained are /10/, /13/:

$$(15a) \qquad \omega_{+}^{n,n+2} = g_c \frac{g_c\left[E - E_{Pauli}(p+1,h+1)\right]^2}{2(n+1)}$$

$$(15b) \qquad \omega_{-}^{n,n-2} = g_c \frac{ph(n-2)}{2} \quad X$$

$$X \left[1 - \frac{(n-1)(p-1)(p-2) + (h-1)(h-2)}{(n-2)8g_c\left[E-E_{Pauli}(p,h)\right]}\right]$$

In (15a,b) the single-particle level density of the compound system is denoted by g_c.

But by the way of the same considerations which have been applied in connection with equations (9) and (10) we can find that the Pauli correction terms in (15a,b) can be neglected as in e q.(9) for excitation energies and particle-hole numbers for which the above formalism is mostly discussed here. Thus we do not present here E_{Pauli} of (15a,b) in detail and refer to the papers /10/ and /13/ which present wrong results corresponding to the differences between (9a) and (9b).

Now in order to obtain a complete theory it would be necessary to calculate $\overline{|M|^2}$. But up to now nobody ever has calculated $\overline{|M|^2}$ in a direct way from a microscopic nuclear model. As an alternative Kalbach /14/ has attempted to find an empirical law for $\overline{|M|^2}$. As such a law Kalbach /14/ made the following proposal

$$(16) \qquad \overline{|M|^2} = K \cdot A^{-3} \cdot E^{-1}$$

hoping that only one universal constant K would be necessary to reproduce the particle emission cross-sections for a wide range of nuclei and excitation energies E.

The above mentioned formalism with relation (16) has been used by the following groups:

Kalbach /10/, Holub, Počanić, Čaplar and Cindro /13/. Fu /15/, Akkermans, Gruppelaar and Reffo /16/ and Gruppelaar, Costa, Nierop and Akkermans /17/.

The STAPRE-code formalism of Strohmaier and Uhl /18/ works with a pre-compound- and compound description separated from the beginning which is not explicitely derived from a common master equation as shown by equations (2)-(9). But the relation (16) is explicitly used in the pre-compound description. Unfortunately K-values have not always been reported by Strohmaier and Uhl. But they do report that K has been used by them as an adjustable parameter.

The intercomparison between /10/, /13/ and /15/ - /17/ is quite problematic because either pairing energy corrections or emitted particle type weighting or both have been taken into account in very different ways. Unfortunately these ways are not always characterized very thoroughly and clearly in the quoted papers so that important details are difficult to recognize. On the other hand Fu /15/ has demonstrated the enormous influence of the way to take into account the level-density pairing-energy correction. This influence can be so strong that one should conclude that this is another source of arbitrariness in addition to the K-problem of (16). Thus we only can intercompare the results respective within each of the papers /10/, /13/ and /15/ - /17/.

In the papers /13/ and /15/ - /17/ the theory is compared with measured cross-sections for 14.6 MeV incident neutrons. Only in the paper /10/ measured charged-particle cross-sections are discussed for incident proton and α-particle energies from 14,6 MeV - 62 MeV within a certain range of the periodic table. As a result of these papers the relation (16) has been roughly confirmed for incident energies from 14,6 - 62 MeV over a range from A = 75 - 200. But because of the different handling of the incorporation of pairing-energy correction and emitted particle type weighting we obtain different K-values for the different papers, namely:

Table 1

paper	K $(MeV)^3$	g $(MeV)^{-1}$	$Q_j(p)$ of eq.(11)
/10/	400	$\frac{A}{13}$	see /10/,/13/
/13/	700	$(\frac{6}{\pi^2})a_{GC}$	see /10/, /13/
/15/	700	$(\frac{6}{\pi^2})a_{GC}$	in initial cond.
/16/	190	$g_c = \frac{A}{13}$; $g_r = (\frac{6}{\pi^2})a_{GC}$ BS	1
/17/	500	$(\frac{6}{\pi^2})a_{BS}$	see /10/,/13/
/36/	400	$\frac{A}{13}$	1

GC = Gilbert + Cameron /19/, BS = Back-Shifted Fermi-gas

r = residual Nucleus , c = compound nucleus

Thus from the above considerations we can conclude that the relation (16) is roughly confirmed for incident energies from 14,6 up to 62 MeV but with the different values of K which are written down above. The preceding formulations with (6) - (16) have been incorporated by F.M. Mann into his computer code HAUSER*5 /20/ where the level density treatment is most similar to /13/. The same is true for the multireaction code GNASH of Young and Arthur /21/. However, for the sake of better mutual comparison the results of work like /10/, /13/ - /17/ and /36/ - especially with respect to the validity of (16) and the value of K - it would be useful to reach a much more uniform description with respect to the level densities (especially concerning the pairing energy- and single-particle level density treatment) as well as to the emitted particle weighting than it is realized in the present situation represented by Tab. 1.

Ambiguities from unsolved level density problems

The differences of the values of K as shown in Tab. 1 for the different publications /10/, /13/ and /15/ - /17/ are as already mentioned partly related to a different handling of the pairing-energy corrections of the compound as well as precompound level densities (exciton state level densities). Thus in the publication /10/ the level density expression (13) was used by C. Kalbach but

with U replaced by $U' = U - \delta$, where δ is the pairing energy correction taken from Gilbert and Cameron /19/. A corresponding pairing energy correction was introduced in the exciton state densities (see (6), (8), (8a)-(8c), (9) and (11)). But the way this has been done is not shown very explicitly in publication /10/. In /10/ $g = \frac{A}{13}$ was chosen as in /16/, (see Table 1).

Contrary to /10/ the Gilbert-Cameron formula /19/ was used instead of (13) in the work of Holub, Počanič, Čaplar and Cindro /13/.
In this work /13/ no pairing-energy corrections were introduced into the exciton state densities of the precompound part because odd-odd compound nuclei or compound nuclei with odd number of the incident nucleon type were investigated. Moreover $g = \frac{6}{\pi^2} a$ was used throughly in /13/ with a taken from Gilbert and Cameron /19/ (see Tab.1) with the corresponding shell effects. But shell effects were also found in /13/ for the K-values of nuclei near closed shells. Here K very much exceeds the average value $K = 700\,(MeV)^3$ (see Tab.1) such as $K = 7000\,(MeV)^3$ for ^{209}Bi and $K = 1400\,(MeV)^3$ for ^{89}Y. But for other nuclei discussed on /13/ such as ^{197}Au with $K = 3500\,(MeV)^3$ and ^{103}Rh with $K = 175\,(MeV)^3$ these K-departures from $K = 700\,(MeV)^3$ cannot so easily be explained as shell effects.

Now Fu /15/ very much stressed that a certain amount of pairing energy always must be expended if a particle-hole pair excitation is accompanied by a pair breaking. Thus pairing-energy corrections must always be taken into account in the exciton state density expressions. But no rigorous derivation of this influence was given by Fu /15/ and consequently no unique results could be obtained. Yet by way of an estimate Fu /15/ could show the strength of this influence. Thus by taking into account this estimate of Fu /15/ the K-value had to be changed from $K = 400\ (MeV)^3$ to $K = 700\ (MeV)^3$. This shows that a rigorous treatment of pairing in the level density expressions of the precompound and compound part with unique results is badly needed in order to give the above formalism a predictive capability with K being not only a fit parameter but a universal constant. This consideration of Fu shows the importance of considering the level densities not only isolated but also in the framework of a consistent nuclear reaction processes. The occurrence of one and the same K-value in /13/ and /15/ although pairing corrections are considered in /15/ and not in /13/ should not be a surprise because the emitted particle weighting is treated correspondingly different in both cases (see last column in Table 1).

In all previously mentioned work equidistant one-particle levels were assumed. The influence of non-equidistance was investigated by Blann and Albrecht /20/ and by Kalbach /21/.

Transition rates from nucleon-nucleon scattering in nuclear matter

Blann /24/and Braga-Marcazzan, Gadioli-Erba, Milazzo-Colli and Sona /25/ went ahead to remove the adjustable parameter K in (14) and (16) by calculating the transition rates $\lambda_{\pm}^{n\pm2}$ in eq.(2) from nucleon-nucleon scattering in nuclear matter according to

$$\lambda_{+}^{1,3} = v\rho\langle\sigma\rangle \tag{17}$$

where v is the particle velocity in nuclear matter

$$v = \sqrt{\frac{2(E+E_F)}{m}}, \tag{18}$$

ρ is the nuclear matter density and $\langle\sigma\rangle$ the effective cross-section for an excited nucleon to interact with nucleons having a Fermi gas momentum distribution. The average $\langle\ \rangle$ is taken over the free nucleon-nucleon scattering cross section with a method due to Goldberger /26/ and Hayakawa, Kawai and Kikuchi /27/ with the Pauli principle taken into account.
The general transition rates $\lambda_{+}^{n,n\pm2}$ then were calculated by Gadioli, Gadioli-Erba, and Sona /28/ using a recursion procedure derived from the expressions (14) and (9). The transition rates thus calculated were then used by Gadioli, Gadioli-Erba, Sona, Sajo-Bohus, Tagliaferri and Hogan /29/ and /30/ in an extended effort to reproduce absolute values of excitation cross-sections for a wide range of mass numbers ($89 \leq A \leq 169$) and excitation energies ($10\ \text{MeV} < E < 100\ \text{MeV}$). But the mentioned authors found they had to multiply the calculated transition rates by factors of the order of 0.1 to 0.25 in order to obtain satisfactory agreement between the calculated and measured cross-sections.
Nevertheless C.Kalbach /33/ has attempted to integrate the more detailed physical knowledge resulting from /29/ and /30/ into an empirical formulation of type (14)-(16) with the result

$$\overline{|M|^2} = \frac{K'}{A^3 e}\left(\frac{e}{7\ \text{MeV}}\right)^{1/2}\left(\frac{e}{2\ \text{MeV}}\right)^{1/2} \qquad e < 2\ \text{MeV} \tag{19}$$

$$= \frac{K'}{A^3 e}\left(\frac{e}{7\ \text{MeV}}\right)^{1/2} \qquad 2\ \text{MeV} \leq e < 7\ \text{MeV}$$

$$|\overline{M}|^2 = \frac{K'}{A^3 e} \qquad 7 \text{ MeV} \leq e \leq 15 \text{ MeV}$$

$$= \frac{K'}{A^3 e} \left(\frac{15 \text{ MeV}}{e} \right)^{1/2} \qquad 15 \text{ MeV} < e;\ e = \frac{E}{n}$$

With (19) and the choice K' = 135 C. Kalbach /31/ was able to reproduce the measured secondary-energy-dependent (p,p') cross-sections of Bertrand and Peelle /32/ for ^{54}Fe and ^{197}Au with incident energies of 29 and 62 MeV in the intermediate secondary energy range. But the high secondary-energy tail came out much too low compared to the measured results of /32/. Nevertheless (14),(15) and (19) have been incorporated by C.Kalbach into her code PRECO-B /33/. Quite good reproductions of experimental results by means of calculations on the basis of (14),(15) and (19) have on the other hand been obtained for (n,2n) and (n,3n) excitation cross-sections by Jhingan, Anand, Gupta and Mehta /34/ for incident energies up to 28 MeV in the mass range 89 to 238. But these cross-sections are not very sensitive to $|\overline{M}|^2$. Gudima, Osokov and Toneev/35/ did not need to reduce λ_+. These authors replaced $E+E_F$ in $\langle\sigma\rangle$ and v by the relative kinetic energy

(20) $$T_n = \frac{8}{5} E_F + \frac{E}{n}$$

of the colliding particles in nuclear matter with n excitons and excitation energy E. Eq.(20) results from the so-called right-angle approximation. T_n is the sum of the mean kinetic energy of an excited particle (p)

(20a) $$T_n^{(p)} = E_F + E_n$$

and the kinetic energy of an intranuclear nucleon (N) averaged over the Fermi spectrum,

(20b) $$T_n^{(N)} = \frac{3}{5} E_F.$$

Gudima, Osokov and Toneev/35/ achieved a good reproduction of the absolute values of the secondary-energy-dependent cross-sections for the reactions Ta(n,n') at 14.6 MeV, Cu(α,p)Zn at 43 MeV and Ta(p,n) cross sections at 14 MeV were calculated in the same way by Hermsdorf, Meister, Seeliger,

Sassonov and Seidel /36/ in good agreement with experimental results on the mass range $30 \leq A \leq 200$.
The absorption cross section σ in Eqs.(6)-(8) was obtained from the optical model. No additional fit parameters were needed but a λ_o-term was added to the master equation with

$$\lambda_o^{n,n} = \frac{2\pi}{\hbar} \overline{|M|^2} g^2 E \quad \frac{3n-2}{2} \quad . \tag{20c}$$

and treated as λ_+ in (17), (18), (20)-(20b).
Tests for more incident energies below as well as above 14 MeV and additional secondary-energy-dependent cross sections for a wide range of mass numbers should be performed before the predictive power of the method can be judged conclusively. This seems necessary in particular because the approximations (17)-(20b) were originally derived for kinetic energies of the colliding particles above about 100 MeV, which means for incident neutron energies above about 55 MeV if we consider $E+E_F$ as a measure for the relative energy of the colliding particles. The applications just mentioned, on the other hand, were made for incident neutron energies well below 55 MeV.

Hybrid and geometry-dependent hybrid model

Blann /24/, /37/, /38/ found out that no fit parameters other than those from the optical model were needed if the excitation energy E in (17) and (18) was replaced by the energy ε of the emitted nucleon, and the Fermi energy E_F by the optical potential depth V. The λ_+ produced this way is then taken the same for each n and is thus independent of n. According to Blann, Kikuchi and Kawai /27/, /37/, /38/ λ_+ can be expressed either by N-N cross section or by optical potential as:

$$\lambda_{j+}(\varepsilon_j) = \begin{cases} \lambda_{j+}^{NN}(\varepsilon_j) & \text{or} \\ \dfrac{2W_j(\varepsilon_j)}{\hbar} = \lambda_{j+}^{W}(\varepsilon_j) \end{cases} \tag{21}$$

where W_j is taken from the imaginary part of the optical potential fitted in the elastic channel of the emitted nucleon. The hybrid model was then obtained by Blann /24/ by inserting (21) into the closed form expression which arises by combining (6)-(8c') after replacing $\lambda_+^{n,n+2}$ by (21) and L(n,E) by the factor before $Q_j(p)$ in the expression (11) for $W_j(n,E,\varepsilon_j)$ divided by the one nucleon level density g_j of a nucleon of type j. This factor is called $\lambda_j(\varepsilon_j)$ according to

$$\lambda_j(\varepsilon_j) = \frac{(2s_j + 1)}{\pi^2 \hbar^3} \frac{\mu_j \varepsilon_j \sigma_j}{g_j} \tag{22}$$

where g_j has to be taken as

$$g_n = (A-Z)/14; \quad g_p = Z/14 \qquad (\mathrm{MeV})^{-1} \tag{22a}$$

Moreover $\lambda_-^{n,n-2}$ is omitted and L(n,E) in all the expressions of (6)-(8c) is replaced by $\lambda_j(\varepsilon_j)$ of eq.(22).
Finally $Q_j(p)$ in (11) is replaced by

$$f_{ij} = \frac{p_{ij}}{p} \tag{23}$$

where p is the total number of particles, p_{ij} the number of particles of type j and t_{ij} the corresponding fraction, given an incident particle of type i. Following Blann /37/ p_{ij} should be calculated according to

$$p_{ij} = \delta_{ij} + \frac{(p-1)\sigma_{ij}}{\sum_{j'} \sigma_{ij'}} \tag{24}$$

where σ_{ij} are the free nucleon-nucleon scattering cross-sections used in a representation which is given in /37/. After the changes introduced with (21) - (24) into (6) - (8c) the question arises whether (6) - (8c) then still can be considered as an approximation of the master equation (2). Blann outlined /37/, /38/ that these changes were suggested to him by considering the formalism of Harp, Miller and Berne /39/, /40/. Because of this composition from two different formalisms Blann calls the resulting formulation the "hybrid model". The resulting expression of the hybrid model for precompound reactions thus becomes after introducting (21) - (24) into (6) - (8c).

$$\left(\frac{d\sigma_{ij}(\varepsilon_i,\varepsilon_j)}{d\varepsilon_j}\right)_{Pr} = \sigma_{ci}(\varepsilon_i) \sum_{n=n_o}^{\bar{n}} \Delta n=2 \; f_{ij} \frac{\omega_{n-1}(U) g_j}{\omega_n(E)} \frac{\lambda_j(\varepsilon_j)}{\lambda_j(\varepsilon_j)+\lambda_{j+}(\varepsilon_j)} D_n \tag{25}$$

$$= \sigma_{ci}(\varepsilon_j) \sum_{n=n_o}^{\bar{n}} \Delta n=2 \; P_{nij}(\varepsilon_j)$$

with

$$D_n = \prod_{n_o+2<n'<n} \left(1 - \sum_{i,j} \left(\int^{E-B_j} P_{nij}(\varepsilon_j) d\varepsilon_j\right)\right) \tag{26}$$

where $P_{nij}(\varepsilon_j)$ is the expression behind the summation sign of (25) and where $\bar{n}$ is the average exciton number at equilibrium obtained from

(27) $$\lambda_{+}^{\bar{n},\bar{n}+2} = \lambda_{-}^{\bar{n},\bar{n}-2}$$

according to (15a,b). The result without Pauli correction is

(28) $$\bar{n} = \sqrt{2gE}$$

From (25) - (28) quite satisfactory results were obtained /41/ for parameter free prediction of secondary energy dependent (α,p) cross-sections for nuclides from ^{51}V to ^{197}Au at 55 MeV incident energy. Only the optical model parameters from the elastic α- and p-channel fits were used and no $|M|^2$-type parameter such as that occuring in (14) and (16) was needed.

Much less successful, on the other hand, were attempts to reproduce the measured angle-integrated secondary-energy-dependent $^{197}Au(p,p')$ cross-section by means of hybrid calculations /42/,/49/. In particular the very flat secondary energy dependence of the measured $^{197}Au(p,p')$ cross-section could not be reproduced by results obtained from calculations on the basis of (25) - (28). These calculated results show a much too steep descent with increasing secondary energy if $n_o = 3$ is chosen. Improvements could be obtained by choosing $n_o = 2$ instead of $n_o = 3$ as the smallest exciton number n_o. But the choice $n_o = 2$ appears quite unphysical unless we assume that at the nuclear surface one of the three initial excitons (a hole) is suppressed. Such an assumption can be understood in the framework of the Thomas-Fermi model, if the Fermi energy, as in the atomic case,is taken as decreasing with the nuclear density d(r) towards the surface according to

(29) $$E_F(r) = \frac{\hbar^2}{2m} \{ \frac{3}{2}\pi^2 d(r) \}^{2/3}$$

where the density follows the Fermi (or Woods-Saxon) distribution

(30) $$d(r) = d_s \{ e^{(r-c)/z} + 1 \}^{-1}$$

with the nuclear half-density radius

(30a) $$c = c_o A^{1/3} , \qquad c_o = 1.07 \text{ fm},$$

the surface thickness

(30b) $$z = 0.55 \text{ fm}$$

and the saturation density

(30c) $$d_s = \{ \frac{4\pi}{3} c_o^3 \}^{-1}.$$

A reasonable way to account for the influence of the nuclear surface diffuseness can be obtained according to Blann /37/ by averaging along the particle trajectory taking the impact parameter

(31) $$R_\ell = \ell\lambda\!\!\!^{-}$$

as the lower limit and the upper limit as

(32) $$R_s = c + 5z$$

outside the nucleus where the density is practically zero. The quantities ℓ and $\lambda\!\!\!^{-}$ in (31) are the orbital angular momentum quantum number and the de Broglie wave length,

(33) $$\lambda\!\!\!^{-} = \frac{\hbar}{\sqrt{2m\varepsilon_o}}$$

The averaged density is then defined by

(34) $$\bar{d}(R_\ell) = \frac{1}{R_s - R_\ell} \int_{R_\ell}^{R_s} d(r)dr.$$

Inserting this into the Fermi energy expression (29) one gets the geometry-dependent Fermi energy or potential depth

(35) $$E_F(R_\ell) = E_F \cdot \{ \frac{\bar{d}(R_\ell)}{d_s} \}^{2/3}$$

where

(36) $$E_F = \frac{\hbar^2}{2m} (\frac{3\pi^2}{2} d_s)^{2/3}$$

is the usual Fermi energy.

From the good results obtained without surface diffuseness for (α,p) reactions by Mignerey and Blann /38/ and Chevarier et al. /41/ with $n_o = 4$ or 5 and from the failure with $n_o = 3$ in the case of $^{197}Au(p,p')$ Blann /43/ concluded that only for $n_o = 3$ (incident nucleons) must the surface diffuseness be taken into account because only then can an exciton acquire enough energy to sense the bottom of the potential well. In this way Blann /43/ found

$$\omega_{1p1h} = gE_F(R_\ell) \,, \qquad U > E_F(R_\ell); \tag{37}$$

$$\omega_{2p1h} = \tfrac{1}{4}g^2E_F(R_\ell)\{2E-E_F(R_\ell)\} \qquad E > E_F(R_\ell).$$

The Ericson or Williams formula (9) is used in all other cases.

In addition there is an influence of the surface diffuseness on the third factor in each sum term of Eq.(25) : g in the expression (22) for $\lambda_j(\varepsilon_j)$ has to be taken as

$$g_j(R_\ell) = \left(\frac{\varepsilon+B+E_F(R_\ell)}{E_F}\right)^{1/2} g_j \tag{38}$$

instead of (22a). Finally also the absorption and excitation rate $\lambda_{j+}(\varepsilon)$ in the third factor of Eq.(25) can be affected by the surface diffuseness. This is the case if $\lambda_{j+}(\varepsilon_j)$ is calculated from the imaginary part $W_j(r)$ of the optical potential for nucleon scattering according to

$$\lambda^{(\ell)}_{j+}(\varepsilon) = \frac{2\bar{W}_j(R_\ell)}{h} \quad \text{with} \quad \bar{W}_j(R_\ell) = \frac{1}{R_s-R_\ell}\int_{R_\ell}^{R_s} W_j(r)dr. \tag{39}$$

In (39) R_s is given by $R_s = r_oA^{1/3} + 5a$ with $r_o = 1.32$ fm and $a = 0.51\,\text{fm} + 0.7\,\text{fm}\,(N-Z)/A$ which is somewhat different from (32).

One can now calculate the pre-equilibrium component of the inelastic-scattering neutron cross section, integrated over emission angles but dependent on the secondary energy, by means of the Eq.(22) - (39). These equations represent the hybrid model with surface diffuseness which was called by Blann the <u>geometry-dependent hybrid model</u>. Apart from general nuclear parameters such as nucleon numbers (N,Z,A), nuclear radius and surface thickness the model contains only the optical-model quantities W and $\sigma(\varepsilon)$. In particular there are no additional fit parameters. Moreover, the geometry-dependent hybrid model is the <u>only existing model that takes the diffuseness of the nuclear surface into account</u>.

On this basis 14.6 MeV (n,n') cross-sections for ^{52}Cr, ^{55}Mn, ^{56}Fe, ^{58}Ni and ^{93}Nb were calculated by Broeders, Broeders and Jahn /44/ (secondary-energy dependent and angle integrated) which are in rather good agreement with the measured results of the groups in Dresden /45/ and Livermore /46/. Also 62 and 39 MeV (p,p') cross-sections of the same kind on ^{56}Fe and ^{209}Bi were calculated in the same way by Blann /38/ who could obtain satisfactory agreement with the measured results of Bertrand and Peelle /32/, and Scobel, Bissem, Friese, Krause, Langanke, Langkau, Plischke, Scherwinski and Wien /47/ compared their identically calculated 27 MeV-(p,p')-results with their own measured results on ^{58,60,61,62,64} Ni and $^{63,65}Cu$ where also good agreement was obtained.

The angular distribution of the secondary-energy-dependent neutrons of the (n,n')-reaction was also introduced into the geometry-dependent hybrid model by the authors of /44/ (see /44/, /49/ and appendix A1) on a PWBA basis. No free fit-parameter was left if the angle integrated secondary-energy-dependent cross-section was equated with that of the n_o=3-component of the geometry-dependent hybrid model. The very satisfactory results obtained for the 14,6 MeV (n,n')-cross-section on ^{56}Fe are shown in the appendix.

Also shown in the appendix is the closure of the gap of measurements between 7,54 MeV and 14,6 MeV incident neutron energy enabled by the geometry-dependent hybrid model because of the absence of any fit parameters other than those of the optical models.

A computer code was developed by Blann /48/ on the basis of this model the first version of which was called ALICE /48/. In this code, as in Refs. /37/ and /43/, the expression

$$g(R_\ell) = \frac{3A}{2E_F(R_\ell)} \qquad (40)$$

was used instead of (38). This led to unrealistic results as described in Ref./49/. The calculations of Hansen, Grimes, Howerton and Anderson (see Ref./50/) were apparently based on Eq.(40) and therefore give too small pre-equilibrium components of the secondary-energy-dependent inelastic neutron-scattering cross section. Also our own first (n,n')calculations on ^{56}Fe and ^{238}U with the hybrid-model code /48/ were only successful after re-introduction of a fit parameter /51/.

+) That such a diffraction type of angular distribution using Bessel functions is more adequate for the mentioned cases than a Legendre polynomial expansion was already pointed out by Pearlstein /81/, Jahn, Broeders and Broeders /82/, Irie, Hyakutake, Matoba and Sonoda /83/ and furthermore by Chiang and Hüfner /77/.

This deficiency of ALICE was corrected in the version OVERLAID ALICE /52/ which was successfully applied to (p,p') reactions by Blann (see Ref. /38/) and to d-, He- and ^{4}He-induced reactions by Bisplinghoff, Ernst, Machner, Mayer-Kuckuk and Jahn, Probst, Djaloeis, Davidson and Mayer-Böricke (see Ref. /38/).

Critical summary of the exciton-master-equation approach

Two groups of precompound descriptions and their applications are reviewed in this report. The first group is based on the master-equation (2) with its two different ways of determining the internal transition rates $\lambda_{\pm}^{n,n\pm2}$. One way consists of reducing $\lambda_{\pm}^{n,n\pm2}$ to a universal empirical law with a universal constant K by using the Golden Rule expressions (14)-(16). But the still too small range and number of examples of incident energies as well as the lack of mathematical transparency of the different versions of calculations does not allow a unique conclusion about the universal law and its constant (see Tab.1 and equations (14) - (16)).
One source for this nonuniqueness is the nonexistence of a unique prescription or at least convection for the incorporation of pairing-energy corrections into the exciton state densities (see the explanations around Tab.1 and in the following paragraph).
For reasons of consistency ambiguities are introduced in this way also into the equilibrium state densities. Moreover all the work based on the attempt of the universal $\lambda_{\pm}$-law (14) - (16) is based on equidistant single-neutron levels. Thus because of all the nonuniqueness mentioned above we have the situation that the question of a universal $\lambda_{\pm}$-law is still in a stage of being explored by fitting measured cross-sections rather than of being used to predict them.

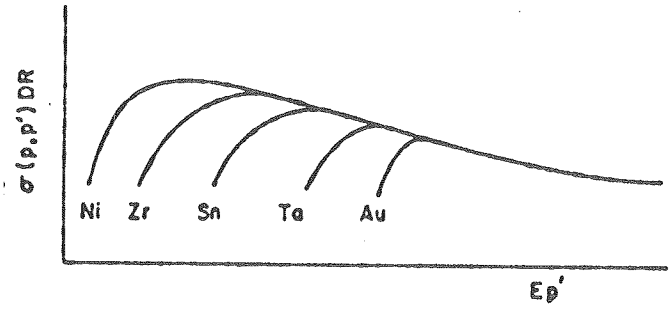
The adherents of the master-equation exciton model approach appear to be very much attracted by its quality of being based on the unique master-equation system, eq.(2), which can be derived directly from the microscopic statistical random matrix model of the nuclear Hamiltonian according to Agassi, Weidenmüller and Mantzouranis /53/. To maintain therefore this exciton master-equation approach the second of the two above-mentioned ways of determining $\lambda_{\pm}^{n,n\pm2}$ was taken by Gadioli et al. /29/,/30/ who went ahead

to fully, calculate the $\lambda_{\pm}^{n,n\pm2}$ -transition rates from nucleon-nucleon-scattering in nuclear matter. But the transition rates resulting from these calculations had to be reduced by 0.1 to 0.25 in order to get full reproduction of the measured (p,x n)-excitation cross-sections for mass numbers $89 \leq A \leq 169$ and excitation energies 10 MeV $\leq E \leq$ 100 MeV. In other words: The calculated cross-sections are too small by factors of 0.1 to 0.25.

As a way out C. Kalbach proposed the still more complicated universal law of eq.(19) trying to reproduce the numerical results of Gadioli et al. /29/, /30/ with the new universal fit-constant K'. But until now this more complicated universal law could be tested with only few examples for only the intermediate secondary energy range and for the less sensitive (n,2n) and (n,3n) cross sections as remarked after eq. (19).
We therefore are inclined to take the result serious that the cross-sections calculated as mentioned above come out too small by 0.1 to 0.25. We think this should be interpreted as an indication that the exciton-master-equation approach does not take into account the full reaction processes. It only takes into account those reactions which are related to the equilibration process. But there should be the direct reaction processes in addition which are not taken fully into account by the exciton-master-equation approach. This is shown very clearly by the formalism of Agassi, Weidenmüller and Mantzouranis /53/ and is also pointed out by Bunakov /54/.

As another strong evidence for the importance of these extra direct reaction contributions it should be considered that the A^{-3}-dependence of $\overline{|M|^2}$ according to equation (16) appears to be empirically rather well established. ω_+ in (14) is predominatly A^3-dependent as shown if the $g \sim A$ behaviour of Tab.1 is introduced into (15a). This means approximate A-independence of λ_+ if the validity of the A^{-3}-dependence of $\overline{|M|^2}$ is assumed according to (16). In other words: The A^3-dependence introduced by (15a) is cancelled by the A^{-3}-dependence of (16). This rises doubts about the Golden Rule treatment of λ_+ as well as about the predominance of the exciton-master-equation contribution at any range of the secondary-energy dependency of the angle-integrated cross-sections. An A-independence of λ_+ gives the results of a slow A-dependence of the angle-integrated secondary energy dependent cross-section according to (6), (8a) and (8b) as about $A^{1/3}$ if integrated over the secondary energy.

But just this type of behaviour is shown by the cross-sections of the direct processes as pointed out by Cohen in the panel of the Albany Conference on Statistical Properties of Nuclei, August 23 - 27, 1971 /73/ who used the following figure:



B.Cohen, Pittsburg

Fig.1:

This figure shows how the strong exponential A-dependent behaviour of the compound-contributions of the (p,p')-emission cross-section at low emission energies goes over to the weakly A-dependent behaviour of the direct contributions as the high-energy tail. Thus this slow A-dependence can be obtained from the direct contributions without the artificial introduction and re-cancellation of the A^3-behaviour shown by the Golden Rule-method to calculate λ_+ of the exciton-master-equation approach.

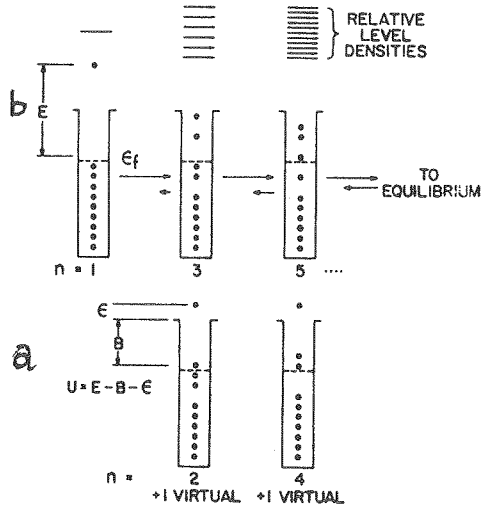


Fig.2: see text

Fig.2a shows a pictorial representation of the first step of a direct reaction process to be taken into account according to /53/, /54/ and /58/-/61/ additionally to the contribution of the exciton master equation approach with its first precompound steps as shown by Fig.2b.

Angular distributions of the exciton-master-equation approach

Further strong indications that no direct reaction processes contribute to the results of the pure exciton-master-equation approach can also be read off from the angular distributions resulting from the angular dependent exciton-master-equation approach developed by Mantzouranis, Weidenmüller and Agassi /55/. These angular distributions show at the high secondary energy tail too small contributions to the forward and backward directions as compared to the measured values /56/, /57/ and /16/.
This can be seen from the results of Mantzouranis /56/ for 45 MeV- (p,n)-reactions on ^{48}Ca, ^{90}Zr, ^{120}Sn and ^{208}Pb, of the author /57/ for 14,6 MeV -(n,n')-reactions on ^{56}Fe and of Akkermans, Gruppelaar and Reffo /16/ for 14.6 MeV-(n,n')-reactions on 33 isotopes from Be till Bi.

The last mentioned results are presented as averaged over secondary energy intervals 2-11 MeV and 6-11 MeV which rises the question whether this means much information in view of the much better resolved secondary-energy spectra measured by the Dresden group /45/ with secondary-energy-bins ranging from 0.5 MeV until 0.05 MeV. Moreover the wide energy-averaging intervals of Akkermans et al. /16/ prevent the pretended possibility of applying their results in the field of fusion reactor design calculations where at least about seven secondary energy groups are needed. A few examples show that smaller energy averaging intervals would much more exhibit the measure of disagreement between the results of these calculations and the experimental results. Finally in order to improve the calculated compared to the measured angular distributions these authors increased artifically and in contrast to the measured behaviour the backward contribution to the N-N cross-section which is an essential input into the method.

We thus consider the foregoing stated deviations between the calculated and the measured angular distributions as a further limitation indicating that the direct reaction processes are not taken into account by the exciton-master-equation approach (see Fig. 2).

Models with explicit account of the direct reaction processes

As a way out it therefore seems to be adequate at first sight to resort to those approaches which take into account the direct processes explicitely in addition to the precompound or compound contributions (as shown by Fig.2).

There are several approaches to calculate direct processes with or without account of precompound or compound contributions which can be grouped according to the names of the following authors:

1. Austern with his book on direct reaction theories /58/.
2. Blokhin, Ignatyuk, Lunev and Pronaev /59/.
3. Tamura, Udagawa and Lenske /60/.
4. Feshbach, Kerman and Koonin /61/.

Approach 1 was developed to treat those direct reaction processes by which single low lying resolved levels of nuclei can be reached. For our context it was used to calculate the high energy tail of the secondary-energy-dependent (p,p')- and (n,n')-cross-sections respectively.
This was done by Fu /62/ for ^{56}Fe on the basis of two DWBA-(p,p') analyses of Peterson /63/ and Mani /64/ of measured cross-sections for 17.5 and 45.35 MeV incident energies. From these (p,p')-analyses in particular of the angular distributions the DWBA-parameter of the first 30 levels were obtained and used by Fu /62/ to calculate the corresponding 14.6 MeV-(n,n')-cross-sections by means of the computer program SALLY /63/. In this way a secondary-energy-distributed 14.6 MeV-(n,n')-cross-section was obtained by Fu /62/ with rather sharp lines around each of the first 25 levels of ^{56}Fe. This DWBA cross-section-distribution obtained for the first 25 discrete levels of ^{56}Fe was then averaged by the author /66/ over the intervals 10-11, 11-12 MeV etc. of the secondary neutron energy and a step-curve was obtained /66/ which agreed quite well with the experimental step-curves of the Livermore /46/ and Dresden /45/ groups.

On the other hand we already mentioned that the smooth curve calculated from the geometry-dependent hybrid model as explained after eq.(38) and presented in /44/ agreed also quite well with the measured Dresden /45/ and Livermore /46/ results. This is in accord with Blanns statements that

he considers the n_o = 3-component of the geometry-dependent hybrid model as the direct component /37/ as demonstrated also by its surface dependence shown by eqs.(29) - (39). This is the only surface dependence shown by any precompound model.

Finally there is the close relationship of the geometry-dependent hybrid model to the Harp-Miller-Berne equations as pointed out by Blann (see the remarks concerning eq. (24)). Now Bunakov /54/ gave a derivation of improved Harp-Miller-Berne equations and showed that the direct contributions are included in them in contrast to the exciton-master-equation approach. The residual interactions of Bunakovs new HMB-equations were completely expressed by the parameters of the optical model /54/. Thus Bunakov's new equations include the direct contributions and do not depend on fit parameters other than those of the optical model. These properties are the same as shown by the hybrid and geometry-dependent hybrid models. Therefore it should be possible to derive those or similar models rigorously from Bunakov's new equations. In case of success we would consider the obtained approach as most preferable against all the other models discussed in this report because the direct contributions would be included and no fit-parameters other than those from the optical model would be needed. But, as already mentioned, this approach 1. till now was only tested for the case of low lying resolved levels.

We therefore have to discuss approaches 2. - 4. developed to calculate the excitation of the unresolved region of levels, the so called continuum part of the spectrum. This was carried out by approach 2. in the random-phase approximation of a phonon model with a self-consistent choice of the effective residual interaction. Up to two phonon excitations were taken into account. Satisfactory reproductions of the measured results were presented for the angular distribution of 20 MeV protons emitted following the impact of 62 MeV protons on ^{54}Fe, of 2-7 MeV secondary neutron from 14 MeV neutron impact on ^{56}Fe for the secondary energy, dependent cross-section of 39 MeV protons incident on ^{54}Fe as well as of 7,9 and 14 MeV neutrons incident on ^{56}Fe. Only rough agreement with the measured results was achieved for the secondary energy dependent cross-sections of 62 MeV protons on ^{54}Fe and on 208 Pb.

Approach 3 has much similarily with approach 2. The only difference is that particle-hole excitations are introduced instead of phonon excitations. Measured angular distributions of 62 MeV incident protons on ^{27}Al and ^{209}Bi are rather well reproduced for secondary proton-energy intervals of 42-52, 32-42 and 22-32 MeV. But Tsai and Bertsch /67/ noted that the energy-weighted sum rule comes out too large with the ph-approximated deformation parameters. Thus Tamura et al. switched to RPA-states and finally to microscopic ph-states and two-step-ph-contributions had to be added /60/. Also Arndt and Reif attempted a similar approach /68/.

Tamura et al. have shown /60/ that approach 4. can be obtained from approach 3. if some simplifications are introduced into the multi-step (predominantly two-step) contributions. So far the one-step contribution of 4. is the same as that of 3. with the difference that the excited level densities are given by RPA response functions in 2. and 3. but by the Ericson ph-function in 4.The latter rises the same level density problems as in the forementioned precompound contributions of the exciton master-equation approach which will become important in particular below 20 MeV excitation energy. Pairing energy corrections are taken into account only in approach 2. The effective interaction of approach 4. has to be adjusted. Good agreement between calculated and measured angular distributions could be obtained by Bonetti et al. /69/ with the same effective interaction force constant V_o = (27,9 $\pm$ 3,5) MeV for angular distributions of 20-40 MeV neutrons emitted from 25-45 MeV protons incident on ^{49}Ca, ^{90}Zr, ^{120}Sn and ^{208}Pb. For the lower incident energies as 25 MeV the statistical multi-step compound contributions of approach 4. become significant, see Bonetti et al. /70/. The appearence of the so called statistical multistep compound contributions in addition to the statistical multi-step direct contributions is a typical aspect of approach 4. which was derived from Feshbachs general framework of nuclear reaction theories /71/ with its P and Q projection-operators onto the open and closed channel spaces leading to both, statistical multi-step direct and statistical multi-step compound contributions. The latter have some similarity with the precompound- and compound contributions of the exciton- master-equation approach. Contributions of this type have not been taken into account by approach 3. for the high-energy examples considered there. A Hauser-Feshbach-contribution has been successfully added

only for the examples of low-energy α-emission cross-sections (< 25 MeV) from 62 MeV protons incident on ^{54}Fe. A similar evaporation contribution has also been taken into account by approach 2.

Whereas in approaches 2. and 3. the effective residual interactions are fixed by self-consistency requirements or sum rules free fit-parameters are left on approach 4. for the residual interactions. Even two strengths of residual interactions were needed in approach 4.: one for the multi-step-direct contributions with V_o = (27 ± 3,5) MeV according to Bonetti et al. /69/ and one for the multi-step-compound contributions with V_o = = 0,70 MeV /70/. But different functions were chosen for the two cases: A Yukawa function for the multi-step direct residual interaction, and a δ-function for the multi-step compound residual interaction. This must be taken into account in comparison of both interaction strengths. But nevertheless they appear to be extremely different, and the question must remain open whether and how this difference can be explained. Moreover this independent adjustability of the two residual interactions of approach 4. can be another source of ambiguity. This has been pointed out by Tamura et al. /60/ by means of the fact that reproductions of measured (p,α)-angular-distributions could be achieved with the same success by a one- step direct plus Hauser-Feshbach approach (see Dragun et al. /72/) as well as by a two-step direct approach (see Tamura et al. /60/). In spite of different incident proton energies in these two cases (44,3 and 34,6 MeV in case of Dragun et al. and 62 MeV in case of Tamura et al.) we consider these two successes with the two different approaches as a hint at the above-mentioned ambiguity which should be investigated somewhat more but which eventually could perhaps be removed by self-consistency requirements or sum rules as in the cases of the approaches 2. and 3.

In any case approaches 1. - 4. demonstrate the occurence of the direct reaction processes and by selecting the advantages it might be possible to obtain a unified and simpler procedure.

Conclusions

The approaches 1. - 3. to take into account the direct reaction processes are substantially able to predict cross-sections whereas approach 4. is a fitting procedure with possible ambiguities. But until now 2. - 4. have been tested only by a few examples. This might have to do with the necessary extensive numerical expense. Simpler is the exciton-master-equation approach. But apparently it does not take into account the direct reaction processes and thus cannot fully describe the forward peaked angular distributions. Moreover it is more a fitting procedure rather than a predictive theory which latter is badly needed to test measured results and to close gaps where measuring is too difficult or even impossible. But as a unique fitting procedure the exciton-master-equation approach could still be useful.
Right now it cannot be obtained this way because of the unnecessary different writing versions of the same solution of the exciton-master-equation approach which is one reason for the different values of the K-constant in Tab.1. Another reason is the lack of a unique procedure to take into account the pairing-energy and shell corrections into the analytic exciton-state and nuclear level density expressions. Also the Pauli-correction term used until now is partially wrong. Thus more consolidation and unification of the very many hitherto existing approaches seem to be necessary rather than still more diversification and blowing up. The mentioned nuclear level density problems occur to a lesser extend if the hybrid and geometry-dependent hybrid models are used. Moreover these models have more predictive capability than the exciton master-equation approach especially for the cases of (n,n')- and (p,p')-reactions (see appendix). But for the calculation of two nucleon and composite particle emission hybrid and geometry-dependent hybrid versions have not yet been developed. Also no collective excitations are included as yet in these models. Moreover no derivation of these models from a basic formalism could be found until now.
It may be remarked that a more preliminary version of this review was already presented by the author as a lecture at the ICTP, Trieste, Jan./Febr. 82 (see /85/).
In a later lecture given by Hodgson at the Varenna-Conference, June 14.-19,1982 (see ref./84/) another review was presented. But in this lecture a different concept was pursued and no attention was paid to the ambiguity problem. On the other hand just this problem has to be considered if progress should be obtained from a mere fitting procedure to a genuine predictive theory.

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APPENDIX

Neutron Emission Cross-Section calculated using Blanns Geometry-Dependent Hybrid Model (GDH).*)

I. Closure of the gap of measurements between 8.56 MeV and 14,6 MeV incident neutron energy.

The GDH is the only precompound model which is free of any fit parameters other than those of the usual optical model. Thus it is the only precompound model which can predict cross-sections and therefore it is the only precompound model which is able to fill the gap of measurements between 8,56 MeV and 14,6 MeV incident neutron energy.

For angle integrated secondary energy dependent neutron emission cross-sections and for the case of ^{56}Fe we have demonstrated this by means of the following diagrams:

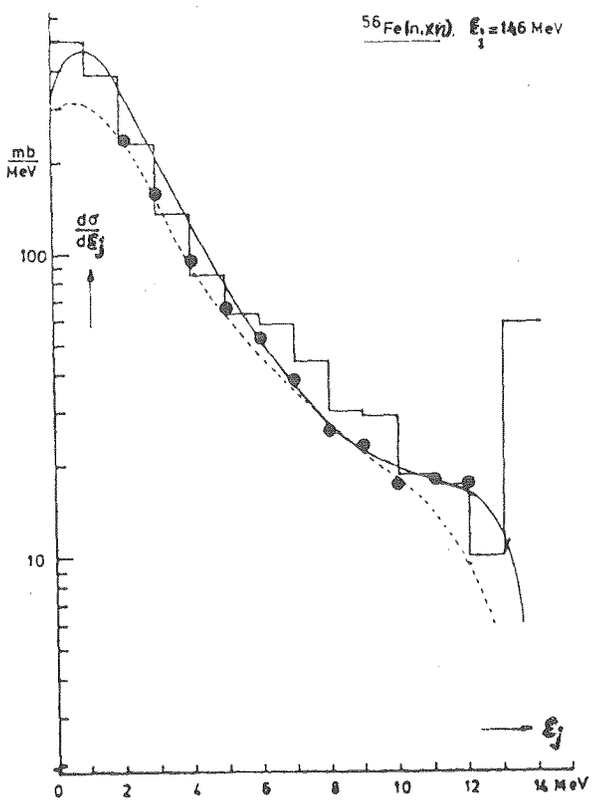
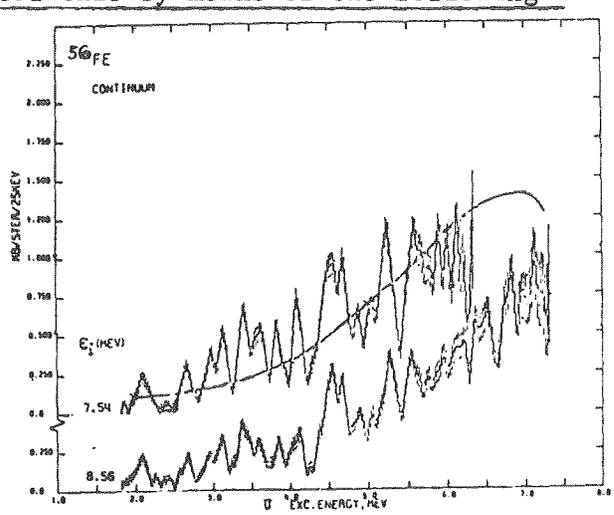


Fig.A1: Results at the upper end of the gap of measurements at incident neutron energy of 14,6 MeV.
Straight line: Our calculated GDH-results /44/ /44/. Points: Dresden measurements /45/.
Step curve: Livermore measurements/46/.
Dotted curve: Calculations from Glauber-theory by Hüfner and Chiang /77/.



$U = \varepsilon_i - \varepsilon_j$ = Excitation energy of the residual nucleus (MeV).

Fig.A2: Results at the lower end of the gap of measurements at incident neutron energy of 7,54 MeV. The fluctuating line represents the measured results of Oak Ridge/74/. Straight line: Our calculated GDH-results./44/ The GDH as a statistical model gives an average through the experimental results.

*) For this first estimate only the one-neutron-channel has been taken into account in the compound-part (Hauser-Feshbach-calculation). Only for ^{93}Nb the full compound (Hauser-Feshbach)-contribution has been considered which was completed by us in /78/.

No other author ever made this test at the both ends of the gap of measurements for the precompound model used by him. But the fulfillement of this test is the necessary precondition for the ability of the used precompound model to fill the gap of measurements.

II. Sucess of GDH for some more nuclei

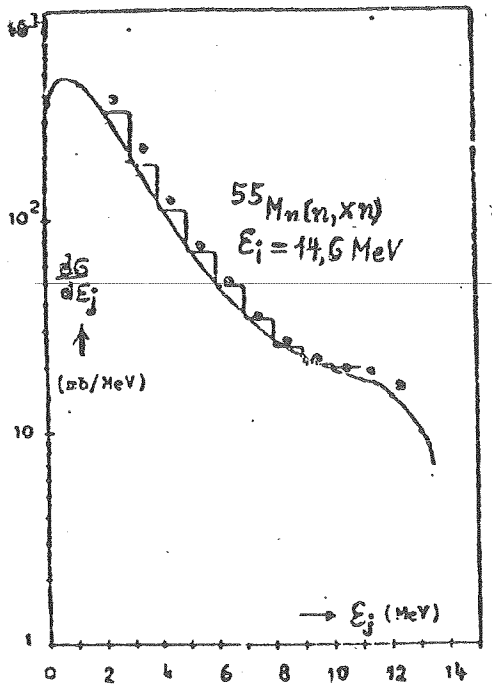


Fig.A3: Angle integrated secondary energy dependent neutron emission cross-section for incident 14,6 MeV neutrons on Mn.
Straight line: Our calculated GDH-results/44/
Points: Dresden measurements /45/
Step curve: Livermore measurements /46/

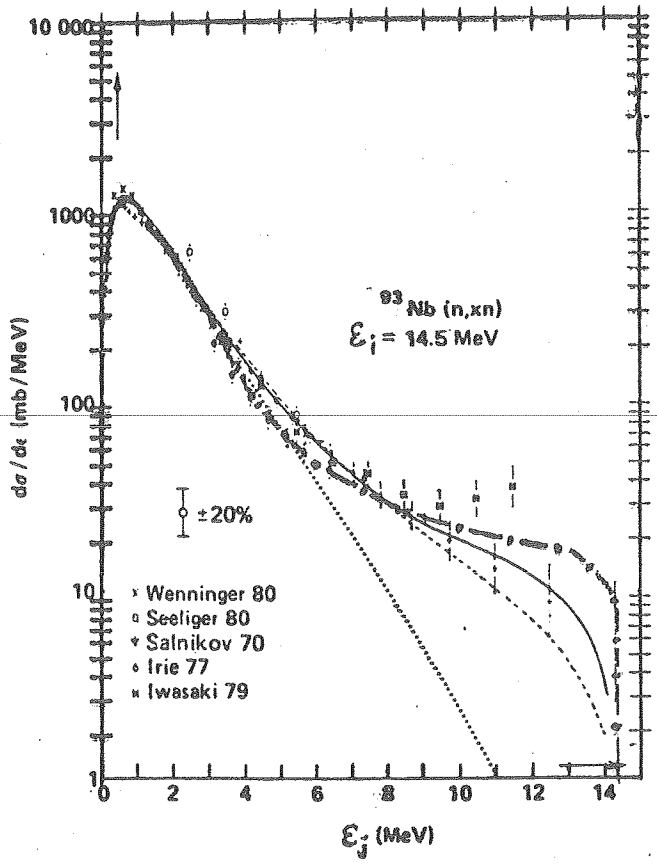


Fig.A4: Angle integrated secondary energy.dependent neutron emission cross-section for incident 14,6 MeV neutrons on Nb.
Straight line:Blanns calculated GDH (N-N-option) + evaporation results/79/
Dashed line:Blanns calculated Hybrid (N-N-option) + evaporation results
Dotted line: " " " "
Dotted-dashed line: Our calculated GDH-results(optical model option)/44/
Experimental Points: (see /79/).

It should be notified once more that all our calculated results shown in Fig. A1-4 are obtained with one and the same GDH with one and the same optical model option without any additional fit parameter other than those of the usal optical model.

Fig.A3 and A4 show angle integrated secondary energy dependent inelastic cross-sections for incident 14,6 MeV neutrons on ^{55}Mn and ^{93}Nb. Our calculated dotted-dashed ^{93}Nb-curve passes as well as Blann's straight curve through the measured points. The high situated points measured at the high energy ^{93}Nb-tail were not shown by the Dresden measurements /46/. But recently such high situated points have been found also by the Dresden group /75/ as direct reaction contributions from collective excitations. They could not yet be presented here.

The corresponding evaluation of the 14,6 MeV neutron cross-section on ^{93}Nb done by the Petten group is shown by the straight line of Fig.A5 as an example for the exciton master-equation approach /80/. The extra fit-parameter K which is characteristic for the exciton master equation approach (see eqn.(16), page 5 and next page 34, III,2 last equation) has been chosen $K = \frac{650A^3}{(13g_r)^3}$ with g_r shown in Table 1. Also Q_j of Tab.1 has been chosen ≠1 with no value presented. A different Q_j-choice seems also to be the reason for the different $K = \frac{500A^3}{(13g_r)^3}$ chosen for a previous ^{93}Nb-fit /17/. Finally a Pb-fit has been obtained with K ten times the first value for $n > n_o$. Again no Q_j-value has been presented.

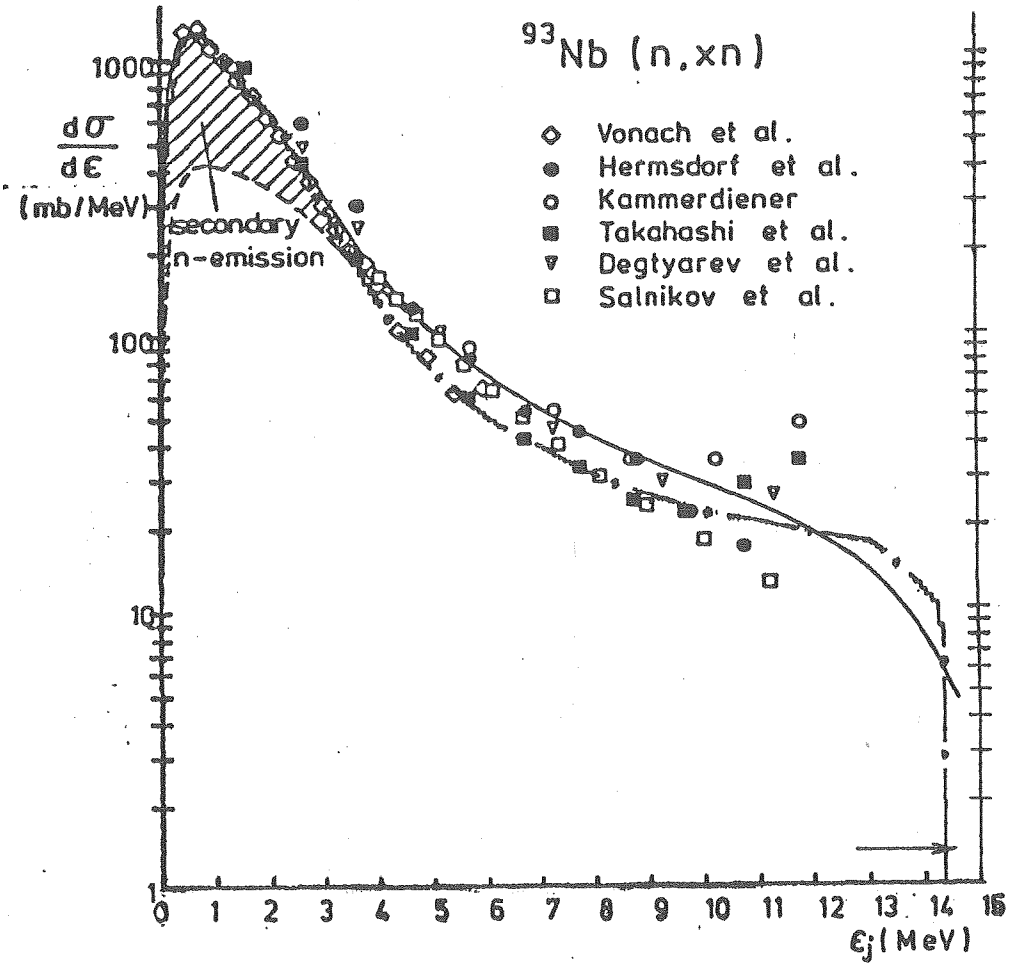


Fig.A5: As Fig.A4 but different model comparison
Straight line: Petten-fit with the exciton master-equation approach /80/
Dotted-dashed-line: Our calculated GDH-results as in Fig.A4 /79/
Experimental points. (see /17, /80/)

III. With or without precompound fit-parameter?

1) GDH - Cross-section (see (21)-(39)):

$$d\sigma_{pr}(\varepsilon_i,\varepsilon_j) = \pi\lambda\!\!\!^{-2} \sum_{\ell=0}^{\infty} (2\ell+1)T_\ell(\varepsilon_i)P_{ij}^{(\ell)}(\varepsilon_j)d\varepsilon_j ; \qquad T_\ell = \text{optical model Transmission coefficient.}$$

$$P_{ij}^{(\ell)}(\varepsilon_j)d\varepsilon_j = \sum_{\substack{n=n_o \\ \Delta_n=2}}^{\bar{n}} f_{ij} \frac{\omega_{n-1}(U)g_j}{\omega_n(E)} \frac{\lambda_j(\varepsilon_i)}{\lambda_j(\varepsilon_i)+\lambda_+^{(\ell)}(\varepsilon_j)} D_n d\varepsilon_j ; \quad n_o=3 \text{ if the incident particle is a nucleon}$$

$$\lambda_+^{(\ell)}(\varepsilon) = \frac{2\bar{W}(R_\ell)}{h} ; \quad \text{with } \bar{W}(R_\ell) = \frac{1}{R_s-R_\ell}\int_{R_\ell}^{R_s} W(r)dr$$

W(r) = Imaginary part of the usual optical potential. Thus no fit parameters other than those of the usual optical model.

2) Exciton Master-Equation Cross-section (EMEC, see (5) - (16)):

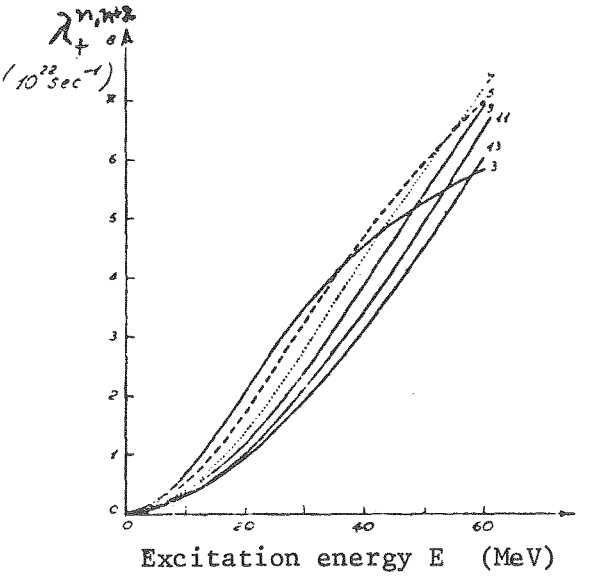
$$d\sigma_{ij}(\varepsilon_i,\varepsilon_j)d\varepsilon = \sigma_{ci}(\varepsilon_i)\sum_{n=n_o} W_j(n,E,\varepsilon_j)T(n,E); \qquad T(n,E) = T^o(n,E)+\alpha w_n$$

$$W_j(n,E,\varepsilon_j) = \frac{2s_j+1}{\pi^2h^3}\mu_j\varepsilon_j\sigma_{ci}Q_j(p)\frac{\omega_{n-1}(U)}{\omega_n(E)}$$

$$T^o(n,E) = \tau(n,E)\left[\prod_{i=n_o}^{n-2,\ n+2}\lambda_+(E)\tau(i,E) + \delta nn_o\right] ; \quad \omega_n = g\frac{(gE - A_{p,n})^{n-1}}{p!h!(n-1)!}$$

$$\tau_n = \frac{1}{\lambda_+^{n,n+2} + \lambda^{n,n-2} + L(n,E)} ; \quad \lambda_+^{n,n+2} = \frac{2\pi}{\hbar}|\bar{M}|^2\frac{g^3E^2}{n+1}$$

$|\bar{M}|^2 = KA^{-3}E^{-1}$ Certainly not valid below 14 MeV (see Fig.A6). Therefore cannot be used to fill gap of measurement.



Gives until more than 4 times too small cross sections.

Fig.A6: Results of calculations of the internal transition rates $\lambda_+^{n,n+2}$ belonging to the exciton master equation cross-section (see III.2, page 34) carried out by Gadioli, Gadioli-Erba and Sona /28/ with the method scetched from eq.(17)-(19). More than 4 times too small cross-sections are obtained and eq.(16) is roughly approximated only above but not below 20 MeV.

IV. Angular Distributions.

Karlsruhe (GDH as in III,1, page 34)

$$(A1) \quad \left(\frac{d^2\sigma_{ij}(\varepsilon_i,\varepsilon_j,\theta_j)}{d\varepsilon_j d\Omega_j}\right)_{dir} = F(\varepsilon_i,\varepsilon_j)\sqrt{\frac{\varepsilon_j}{\varepsilon_i}}\,(2\ell'+1)\sum_L C_{\ell\ell'}(LO,OO)\,j_L^2(QR) \qquad +)$$

$$(A2) \quad \int d\Omega_j \left(\frac{d^2\sigma_{ij}(\varepsilon_i,\varepsilon_j,\theta_j)}{d\varepsilon_j d\Omega_j}\right)_{dir} = \left(\frac{d\sigma_{ij}(\varepsilon_i,\varepsilon_j)}{d\varepsilon_j}\right)_{n=n_o=3} \quad ; \text{ isotropic for } n>3$$

ℓ,ℓ' are ground- and excited state shell model angular momenta where $\ell+\ell' > L > |\ell-\ell'|$, $F(\varepsilon_i,\varepsilon_j)$ is to be calculated from (A2), $j_L(QR)$ are spherical Bessel functions of the first kind, Q momentum transfer, R nuclear radius and $C_{\ell\ell'}$ Clebsch-Gordan coefficients. The results of Fig.A7 are obtained with only one L respectively.

+) see foot-note on page 17.

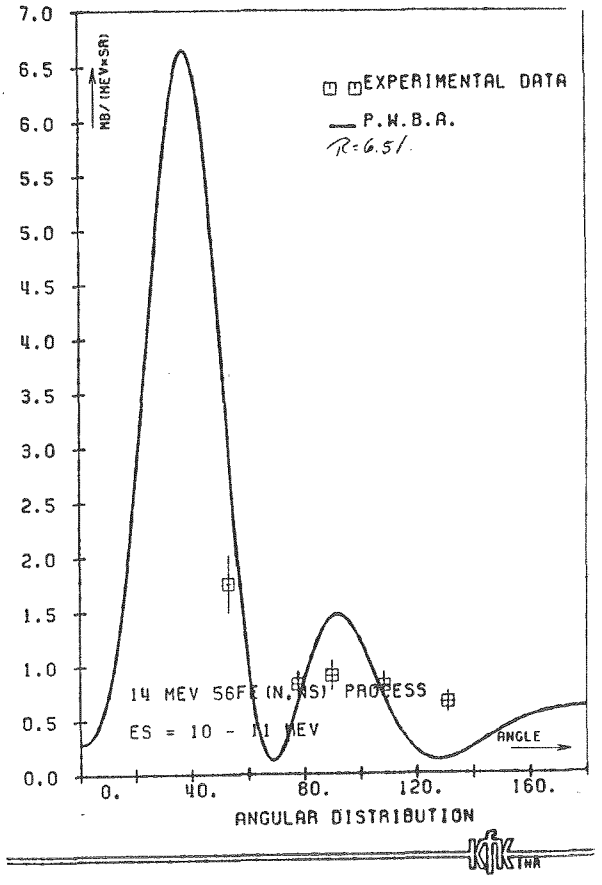


Fig.A7: Angular distributions of secondary energy dependent neutrons emitted after impact of 14.6 MeV neutrons on

a) ^{56}Fe-target, secondary energy interval 10 - 11 MeV

b) ^{93}Nb-target, secondary energy interval 8 - 9 MeV

Points: Measured results of Dresden group /45/.

Curves: Straight curves: Calculated results according to the preceding formulas (A1) and (A2) of page 35 with L=2, R=6,51 fm for ^{56}Fe and L=3, R=7,11 fm for ^{93}Nb (see /44/ and /78/).
Dashed curve in Fig. 7b: Calculated results of the Petten group for a secondary energy of 9 MeV /76/.

Fig.A7 a

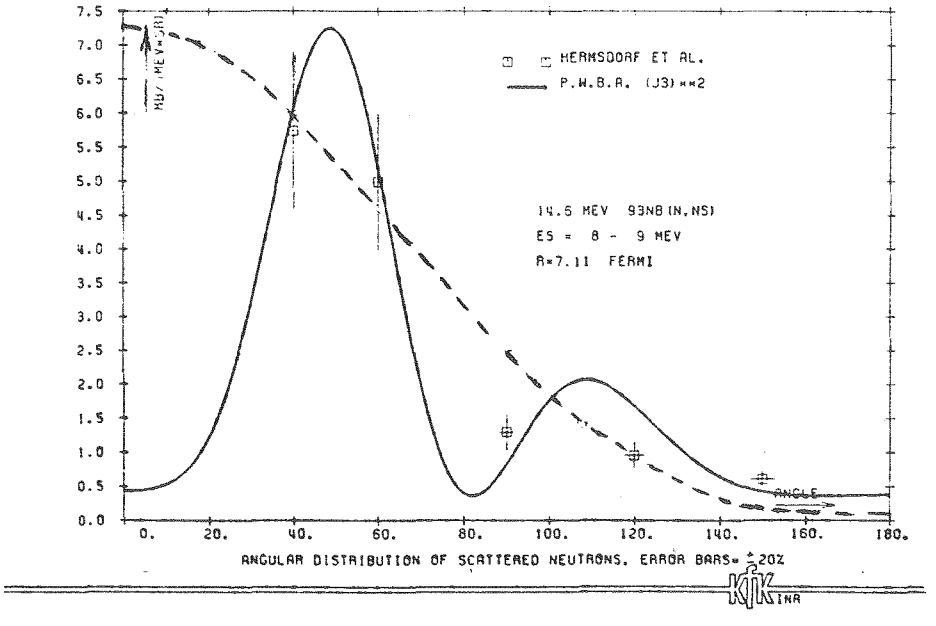


Fig.A7 b

The Bessel-type of angular distribution which is typical for direct processes is shown by our straight curves in Fig. 7 obtained according to (A1) and (A2) while this Bessel-type of behaviour is not shown by the dashed curve of the Petten-group /76/ plotted in Fig. 7b. (see the comments on p. 21 and foot-note on page 17).

V. Survey of Literature

1.) EXCITON-MODEL: GRIFFIN/BLANN, 1966/68

TIME-DEPENDENT APPROACHES

2.) EXCITON-MASTER-EQUATION: CLINE-BLANN, 1971
APPLIED BY: KALBACH, OBLOŽINSKÝ, RIBANSKÝ AND BĚTÁK, UHL-STROHMAIER, YOUNG-ARTHUR, HOLUB-POČANIĆ-ČAPLAR-CINDRO, GRUPPELAAR, GADIOLI ET AL.

3.) HARP-MILLER-BERNE (HMB)-MODEL, 1971
WITH IMPROVEMENT DERIVED BY BUNAKOV, 1976/77

4.) HYBRID-MODEL: BLANN, 1971

5.) GEOMETRY-DEPENDENT HYBRID-MODEL (GDH): BLANN, 1972/73
APPLIED BY: BLANN, JAHN-MACHNER-MAYER-BÖRICKE, BISPLINGHOFF-ERNST-MAYER-KUCKUK, C. & I. BROEDERS-JAHN, SCOBEL ETC.

6.) DERIVATION OF THE MASTER-EQUATION + DIRECT-KOMPONENT: AGASSI, WEIDENMÜLLER & MANTZOURANIS, 1975

7.) MASTER-EQUATION WITH ANGULAR DISTRIBUTION: MANTZOURANIS, WEIDENMÜLLER & AGASSI, 1976
APPLIED WITHOUT DIRECT-KOMPONENT BY: MANTZOURANIS (ONLY (P,N)-PROCESSES), GRUPPELAAR (14,6MEV(N,N'))

DEVELOPMENT

DEVELOPMENT

CHRONOLOGICAL ORDER

TIME-INDEPENDENT APPROACHES

8.) STATISTICAL MULTI-STEP COMPOUND (SMSC) + STATISTICAL MULTI-STEP DIRECT (SMSD): FESHBACH-KERMAN-KOONIN (1980)
APPLIED BY: BONETTI-COLLI-HODGSON (1980/81)

9.) CONTINUUM-DIRECT: TAMURA, UDAGAWA AND LENSKE (1982)

10.) SEMI-MICROSCOPIC-CONTINUUM-DIRECT: BLOKHIN, IGNATYUK, LUNEV, PRONYAEV (1977/79)

11.) NUCLEAR MATTER-DIRECT: V.GERAMB (1982), (see /84/)

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