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A b s t r a c t

The validity of a resonance approximation for the reaction amplitudes is discussed. A method of determination of the dominant mechanism of a reaction ($1/\sqrt{N}$ - classification of amplitudes, N is the number of component in the wave function of a compound state) is suggested.

Now the bulk of information about parity violation in neutron reactions has been stored. Parity violation in the reaction (n,γ) was discovered in work /1/. One can find the references to the works performed prior to 1976 in the review /2/ (see also the more recent works /3-5/). Parity violation in nuclear fission in the reaction (n,f) was discovered in the work /6/. The theory of this phenomenon and the references to the following experimental works are presented in the review /7/ (see also Refs. /8-11/).

In the recent years the parity violation effects which arise when the neutrons propagate through the matter, are extensively studied. The spin rotation of cold neutrons around the direction of motion in ^{117}Sn (Ref. /12/) and the difference in the capture cross sections of the right- and left-polarized neutrons in ^{117}Sn, ^{139}La, ^{81}Br (Refs. /12-14/) have been found. Finally, the difference in the neutron capture cross sections near the p-wave compound resonances where the relative magnitude of the effect constitutes 10^{-1}-10^{-2} has been measured /15-17/. In all the cases listed above, the parity violation effects are enhanced substantially because of the complexity of compound states (dynamical enhancement) and the closeness of a p-wave compound resonance /18/ (see also Refs. /19,7,20,21/). It is clear that a considerable enhancement of the P-odd effects near the p-wave compound resonances occurs in the (n,γ) reaction as well /18,7/. It should be noted that long before publication of these papers, the formula has been presented in the Ref. /22/, which contains the enhancement of the circular polarization of γ-quanta near a p-wave compound resonance. Nevertheless, the existence of a resonance enhancement was not emphasized in Ref. /22/ and the standard estimate $P_\gamma \sim 10^{-4}$ (which coincides with the well-known value at the thermal point) for the magnitude of the circular polarization in the resonance was given.

In connection with the study of the parity nonconserving effects in nuclei, the problems arise which are also important, in our opinion, for a study of the other phenomena in nuclear reactions:

1. What is the range of validity of the resonance appro-

ximation in the (n,n), (n,γ) and (n,f) reactions?

2. How can the dominant mechanism leading to the appearance of the effect. be found?

The present paper is devoted to the discussion of the above problems.

1. Resonance approximation in heavy nuclei

The observed variation of the parity violating effects from a p-wave resonance to the thermal point constitutes 3-4 orders of magnitude. The natural question arises: are the simple resonance formulae valid in such situation? The problem of a resonance approximation in nuclear reactions has long history. It is enough to recall the R-matrix Wigner-Eisenbud theory /23/ or the Kapur and Peierls theory /24/ (see also the book by Lane and Thomas /25/). However, the technique developed in these works is cumbersome and is not convenient, in our opinion, for consideration of the angular and polarization correlations due to a large number of channels connected with the reaction (n,γ). We prefer to use the method based on the summation of the series of perturbation theory. In essence, this method has a great deal in common with the approach suggested by Feshbach /26/.

Let us first consider the elastic neutron scattering in the vicinity of an isolated compound resonance. The wave function of a compound state may be represented as a sum of the products of single-quasiparticle wave functions. Let us exclude the single-particle component from this function (i.e. the states with neutron above the non-excited initial nucleus). A compound state $|c\rangle$ defined by such a way is the level of discrete spectrum because in the vicinity of the neutron threshold the remaining components have no output to the continuous spectrum (the other decay channels are not yet taken into consideration). The emittance and capture of a neutron (i.e. the coupling with a single-particle component) will be taken into account using the perturbation theory in the residual interaction H_s. The scattering amplitude consists of the potential and resonance parts: $f = f_0 + f_2$. An expression for the resonance part is of the form

$$f_2(E) = \frac{\langle n|H_s|c\rangle\langle c|H_s|n\rangle}{E - E_c + i0} + \int \frac{\langle n|H_s|c\rangle\langle c|H_s|n'\rangle\langle n'|H_s|c\rangle\langle c|H_s|n\rangle}{(E-E_c+i0)(E-E'+i0)(E-E_c+i0)}\rho(E')dE' + \dots \quad (1)$$

Here $|n\rangle$, $|n'\rangle$ are the single-particle wave functions, E is the neutron energy, E_c is the energy of the compound state, $\rho(E')$ is the density of single-particle states (phase volume). Of course, not only the integral over dE' but also the summation over single-particle levels of the discrete spectrum are carried out in formula (1). The graphs corresponding to formula (1) are shown in Fig. 1. Emphasize that we sum all the orders of perturbation theory, i.e. we do not assume the smallness of operator H_s. As known, the series (1) is a simple geometric progression. If one introduces the Green's function of a compound nucleus $G_c(E)$ and mass operator $\Sigma(E)$ then f_2 is written as follows:

$$f_2(E) = \langle n|H_s|c\rangle G_c(E)\langle c|H_s|n\rangle \quad (2)$$

$$G_c(E) = \frac{1}{E-E_c+i0}\left[1 + \frac{\Sigma}{E-E_c+i0} + \left(\frac{\Sigma}{E-E_c+i0}\right)^2 + \dots\right] = \left(E - E_c - \Sigma(E)\right)^{-1} \quad (3)$$

$$\Sigma(E) = \int \frac{\langle c|H_s|n'\rangle\langle n'|H_s|c\rangle}{E-E'+i0}\rho(E')dE' = \Delta E_c(E) - i\frac{\Gamma_c(E)}{2} \quad (4)$$

In the case of an arbitrary number of resonances the Green's function and the mass operator become the matrices in space of the compound states:

$$\hat{G} = \left(\hat{G}_0^{-1} - \hat{\Sigma}\right)^{-1}, \quad G^{(0)}_{cc'} = \frac{1}{E - E_c + i0}\delta_{cc'} \quad (5)$$

$$f_2 = \sum_{cc'}\langle n|H_s|c\rangle G_{cc'}(E)\langle c'|H_s|n\rangle \quad (6)$$

The other decay channels (for example, the γ-channel) give the additional terms in the mass operator: $\hat{\Sigma} = \hat{\Sigma}^n + \hat{\Sigma}^\gamma + \dots$

$$\Sigma^n_{cc'} = \int \frac{\langle c|H_s|n'\rangle\langle n'|H_s|c'\rangle}{E - E_c + i0}\rho_n(E')dE' \quad (7)$$

$$\sum_{cc'}^{\gamma} = \sum_{q} \int \frac{\langle c|H_{em}|\gamma q\rangle \langle \gamma q|H_{em}|c'\rangle}{E-\omega-E_q+i0} \rho_{\gamma}(\omega)\, d\omega \qquad (8)$$

Here H_{em} is the operator of electromagnetic interaction, index q enumerates the nuclear states to which the γ-transition takes place. Using the Green's function of a compound nucleus, one can write the amplitude of any process, for example, of the reaction (n, γ):

$$f_{\gamma} = \sum_{cc'} \langle n|H_s|c\rangle \zeta_{cc'} \langle c'|H_{em}|\gamma q\rangle \qquad (9)$$

In order to write f_γ as a simple sum of resonance terms, it is necessary to diagonalize the Green's function $\zeta_{cc'}(E)$. According to formulae (6), this is equivalent to the diagonalization of the matrix $A_{cc'} = E_c\delta_{cc'} + \sum_{cc'}(E)$ (cf. formulae (3.17) and (3.18) from Ref. /26/):

$$\tilde{A} = U^{-1} A U \quad , \quad \tilde{A}_{cc'} = \tilde{E}_c(E)\,\delta_{cc'} \quad ,$$
$$\tilde{\zeta} = U^{-1}\zeta U \quad , \quad \tilde{\zeta}_{cc'} = \frac{1}{E-\tilde{E}_c(E)}\,\delta_{cc'} \quad . \qquad (10)$$

It is known that in general case the matrix U is a nonunitary. As a result of diagonalization, we obtain, therefore, the right- and left-hand sets of the states (see, e.g., Ref. /27/):

$$|\tilde{c}\rangle = \sum_{c'} U_{c'c}\,|c'\rangle$$
$$\langle\tilde{c}| = \sum_{c'} U^{-1}_{cc'}\langle c'| \qquad (11)$$

Generally speaking, the eigenvalues $\tilde{E}_c = Re\,\tilde{E}_c + i\,Im\,\tilde{E}_c$, the matrix U, and the sets of states $|\tilde{c}\rangle$, $\langle\tilde{c}|$ can depend on energy E, that substantially complicates the behaviour of amplitude f_γ. But it turns out that practically there is no this dependence in complex nuclei. The reason is that all typical variation scales for the matrix elements are of the order of several MeV and are large compared to the average distance between the compound resonances $\mathcal{D}$. Consider, for example, the neutron contribution to $\sum$:

$$Re\sum_{cc'}^{n} = \int \frac{\langle c|H_s|n'\rangle\langle n'|H_s|c'\rangle}{E-E'} \rho_n(E')\,dE' \sim \Gamma_n(E_m) = \qquad (12)$$
$$= \left(\frac{\Gamma_n^{(0)}}{\mathcal{D}}\sqrt{\frac{E_m}{T_{ev}}}\right)\cdot\mathcal{D} \lesssim \mathcal{D}$$

Here $\int$ is the principal value of integral, $\Gamma_n^{(0)}$ is the neutron width at E = 1 eV, E_m is the characteristic energy at which the integral converges. The convergence of an integral is determined by an energy dependence of the matrix elements. Hence, $E_m \sim \Gamma_{spr} \sim$ MeV, Γ_{spr} is the spread width of single-particle resonances. In estimation we have taken the data on the strength function for s-levels from Ref. /28/: $\Gamma_n^{(0)}/\mathcal{D} \lesssim 10^{-4}$. For the derivative of $Re\sum$, we obtain (at $\mathcal{D} \sim$ 1-100 eV):

$$\frac{\partial\, Re\sum^{n}}{\partial E} \sim \frac{Re\sum^{n}}{E_m} \lesssim \frac{\mathcal{D}}{E_m} \lesssim 10^{-4} \div 10^{-6} \qquad (13)$$

Thus, at the variation of energy $\delta E \sim \mathcal{D}$, the relative shift of the resonance position does not exceed 10^{-4}. As far the imaginary part, $Im\sum^{n}$ can be neglected near the neutron threshold. The origin of $Re\sum^{\gamma}$ is the same as that of the Lamb shift in atoms. There is no difficulty to verify that $Re\sum^{\gamma} \sim Im\sum^{\gamma} \sim \Gamma_\gamma \lesssim \mathcal{D}$. The characteristic frequencies of the emitted γ-quanta and the scale at which matrix elements of H_{em} vary constitute 1-10 MeV. Hence, $\delta Re\sum^{\gamma}/\mathcal{D} \sim 0.1\mathcal{D}/\text{MeV} < 10^{-5}$ and $\delta Im\sum^{\gamma}/\Gamma < 10^{-4}$. Similar estimates can be presented for fission as well.

As we have mentioned above, the matrix U is not, generally speaking, an unitary one. This circumstance leads to the appearance of additional phases in the matrix elements and to the distotrion of the energy dependence of the total cross section and interference effects (including P-odd ones). Such a situation can occur if the fission channel is open. However, this is not to be so if only the neutron and γ-channels are open and we analyse the reactions near the neutron threshold. In this case, U proves to be an unitary real matrix. In order to verify this, it is appropriate to take the intermediate states in formulae for mass operator (7) and (8) in the form of standing waves. Then, the matrix elements

in these formulae are real and one can immediately see that $Re\Sigma_{cc'}$ is a symmetric matrix. The imaginary part, in this case, turns out to be proportional to the unit matrix: $Im\Sigma_{cc'} = -\frac{\Gamma}{2}\delta_{cc'}$. The neutron contribution to $Im\Sigma$ can be neglected, since $Im\Sigma^n \sim \Gamma_n \ll \Gamma_\gamma$ near the threshold. As known, the fluctuations of the γ-widths are small, $\delta\Gamma_\gamma \ll \Gamma_\gamma$, because of a large number of γ-channels. In addition, the non-diagonal part $Im\Sigma^\gamma_{cc'}$ is small since the statistical relation (see Appendix A) holds:

$$\overline{\left(Im\Sigma^\gamma_{cc'}\right)^2} = \frac{1}{8}\overline{(\delta\Gamma)^2} \qquad (14)$$

Thus, in order to diagonalize the Green's function, it is sufficient to make a diagonalization of the real symmetric matrix $A_{cc'} = E_c\delta_{cc'} - Re\Sigma_{cc'}$, which is performed by a simple rotation in space $|c\rangle$.

As a result, the resonance part of the neutron scattering amplitude near neutron threshold ($\Gamma_n \ll \Gamma_\gamma$) is found to be a simple sum of resonance terms having no additional phases. Their energy dependence is due to the resonance energy denominators and the kinematic dependence of amplitudes (K^ℓ, K is the neutron moment). The same assertions are also valid for the reaction (n,γ). If the fission channel is open, additional phases, $\varphi \sim \delta\Gamma/D$, can arise.

In calculating the parity violating effects one needs to take into account a weak interaction, basing on perturbation theory. From the conclusion drawn above, it is clear that the matrix element of weak interaction in the admixture amplitudes should be calculated using the compound states $|\tilde{c}\rangle$, and the diagonalized Green's function of a compound nucleus should be used. Since in the calculations the parameters of the compound states are extracted from the experiment, the use of the states $|\tilde{c}\rangle$, instead of $|c\rangle$, does not complicate the problem.

In conclusion of this section, we would like to make one simple note. As known, in the one-channel problem the description of an amplitude as a sum of Breit-Wigner resonance terms contradicts to the unitarity condition. This is

easy to see from the optical theorem:

$$\sigma = \frac{4\pi}{k} Im f(0)$$

When squaring the amplitude, the interference terms appear in the left-hand side of the equality whereas the right-hand part contains only the sum of the resonance contributions. In the case of the reactions (n,γ) and (n,n) near the neutron threshold, which has been discussed above, the applicability of simple resonance formulae is due to the fact that the interference terms have no definite sign and are averaged in summation over a variety of final states connected with the reaction (n,γ) (due to a random sign of the interference terms, their relative contribution to the total cross section is $\sim 1/\sqrt{N_f}$, N_f is an effective number of final states giving the main contribution to the cross section (n,γ)). These arguments fail far from the neutron threshold (at $\Gamma_n \gtrsim \Gamma_\gamma$) because there is an interference between the resonances in the neutron part of the cross section.

II. Matrix elements of one-particle operators between the compound states

Let us represent the wave function of a compound state as an expansion in products of single-quasiparticle wave functions

$$\psi = \sum C_\alpha \varphi_\alpha \qquad (15)$$

$\varphi_\alpha = a_i^+ \ldots\ b_k^+ \ldots\ |0\rangle$ are the states which contain a definite number of particles and holes. (a^+,b^+ - creation operators for particles and holes. We are not concerned with the collective degrees of freedom since their inclusion does not change the situation qualitatively).

The characteristic number of the "principal" components in ψ (i.e. the components which give the main contribution to normalization) is equal to (see Refs. /28-30,7/)

$$N \sim \frac{\pi}{2}\frac{\Gamma_{spr}}{D} \gg 1 \qquad (16)$$

Here $\mathcal{D}$ is the average distance between the compound levels with the same momentum and parity, $\Gamma_{sp2} \sim$ MeV is the fragmentation width of the state φ_α to the levels of the compound nucleus. By virtue of the normalization condition for the principal components $|C_\alpha| \sim 1/\sqrt{N}$. The components φ_α, whose unperturbed energy (i.e. the energy without taking into account the residual interaction H_s leading to fragmentation) considerably differs from that of the compound state ($|E_\varphi - E_\psi| > \Gamma_{sp2}/2$), appear in ψ with a smaller weight. The corresponding coefficients are easy to estimate using the perturbation theory in residual interaction H_s:

$$C_\alpha = \frac{\langle \varphi_\alpha | H_s | \psi^{(0)} \rangle}{E_\psi - E_\varphi} \sim \pm \frac{\Gamma_{sp2}}{2(E_\psi - E_\varphi)} \frac{1}{\sqrt{N}} \qquad (17)$$

Here $\psi^{(0)}$ is the wave function of the compound state with the small components not taken into account. We have used the formula from Ref. /28/:

$$\Gamma_{sp2} = \frac{2\pi}{\mathcal{D}} \overline{\langle \varphi_\alpha | H_s | \psi^{(0)} \rangle^2} = \frac{4}{\Gamma_{sp2}} N \overline{\langle \varphi_\alpha | H_s | \psi^{(0)} \rangle^2} \qquad (18)$$

In this book the formula is presented which describes the coefficients for both the principal and small components:

$$\sqrt{\overline{C_\alpha^2}} = \frac{\Gamma_{sp2}/2}{\sqrt{(E_\psi - E_\varphi)^2 + \Gamma_{sp2}^2/4}} \frac{1}{\sqrt{N}} \qquad (19)$$

Formulae (18) and (19) have been derived in Ref. /18/ for the model where the density of the states is constant. In the nucleus, the density of the states increases exponentially with energy because of increasing the number of quasiparticles in the state φ_α. But H_s is a two-body operator, and hence, according to Eq. (17), only states φ_α, which differ from the principal components not more than by the state of two nucleons, admix to $\psi^{(0)}$. The density of such states weakly varies with energy. In view of this, if the total density of the states varies not strongly at $\delta E \sim \Gamma_{sp2}/2$, the estimates (17), (18) and (19) are also applicable to a real nucleus. The difference consists only in the fact that not all the remote components admix to ψ,

and $N = N(E\psi)$ in the formulae.

In the nuclei (Sn, Cd, La, Br, U, Pu) where the parity violating effects in neutron reactions have been observed, $N \sim 10^4 \div 10^6$. The large value of the parameter N which characterizes the complexity of the system, enables one to readily estimate the relative contributions of various mechanisms to the reaction amplitude. As will be shown below, these contributions are proportional to different powers of $\sqrt{N}$.

Let us first estimate the matrix element of a one-body operator $\hat{M}$ between the compound states (Refs. /31,29,30,7/). As $\hat{M}$, one can take the operator of electromagnetic, weak, or Coriolis interaction.

$$M_{if} = \langle \psi_f | \hat{M} | \psi_i \rangle = \langle \sum_\alpha a_\alpha \varphi_\alpha | M | \sum_\beta b_\beta \varphi_\beta \rangle = \qquad (20)$$
$$= \sum_{\alpha\beta} a_\alpha^* b_\beta \langle \varphi_\alpha | \hat{M} | \varphi_\beta \rangle$$

For the sake of definiteness, $N_f \leq N_i$ is assumed. With a fixed α, the matrix element $\langle \varphi_\alpha | \hat{M} | \varphi_\beta \rangle$ is not equal to zero at several β (their number is denoted by q) when φ_β differs from φ_α by the state of one particle only. It is naturally to think that the signs of particular terms in the sum (20) are random in such a complex system as the compound nucleus. Thus we have, in eq. (20), the incoherent sum of $\sim qN$ terms, each of the order of $M/\sqrt{N_i N_f}$, M is typical single-particle matrix element. As a result of the summation of incoherent contributions,

$$\overline{M_{if}^2} \sim q N_f \left(\frac{1}{\sqrt{N_i}} \cdot \frac{1}{\sqrt{N_f}} \cdot M \right)^2, \qquad (21)$$
$$M_{if} \sim \sqrt{q}\, M/\sqrt{N_i}$$

In the estimation performed, the matrix element between the principal components of the states $|i\rangle, |f\rangle$ has been assumed to be different from zero. But, by virtue of the selection rules for operator $\hat{M}$, the matrix elements between the principal components can prove to be zero. In this case, the matrix element is distinct from zero due to the small components in ψ_i, or ψ_f. For example, such a situation occurs if we analyse the matrix element of weak interaction between the

close compound states (Ref. /32/). The operator of weak interaction H_w transfers the particle from one shell to the other. Acting on the principal component of the state $|f\rangle$, H_w transfers it to the component lying in the another shell. According to formula (17), or (19), the coefficient C_α for such a component in the state $|i\rangle$ is suppressed by a factor $\Gamma_{sp}/(2\omega_0)$ ($\omega_0 \sim 8$ MeV) is the distance between the shells). As a result,

$$(H_w)_{if} = \sum_{\alpha\beta} \left[a^*_\alpha(small) \cdot b_\beta \langle\alpha|H_w|\beta\rangle + a^*_\alpha b_\beta(small) \langle\alpha|H_w|\beta\rangle \right] \sim \tag{22}$$

$$\sim \frac{\mathcal{D}}{\omega_0} H_w \sqrt{q N_f} \sim \sqrt{q}\, \frac{\Gamma_{sp}}{\omega_0} H_w / \sqrt{N_i}$$

Since Γ_{sp} depends, generally speaking, on the number n of quasiparticles in the state φ_α, formula (22) includes the effective width corresponding to the states with the most probable n. As a result eq. (22) contains in comparison with the estimate (21) the factor of suppression $\Gamma_{sp}/\omega_0 \sim \frac{1}{5} \div \frac{1}{10}$.

Using Eq. (22), it is easy to estimate the mixing coefficient of the nearest compound states with different parity:

$$\alpha \sim \frac{(H_w)_{if}}{\mathcal{D}} \sim \frac{H_w}{\omega_0} \sqrt{Nq} \tag{23}$$

The ratio H_w/ω_0 is a typical magnitude of mixing between single-particle levels. The mixing of compound levels is enhanced by $\sqrt{qN}$ times. This corresponds to the known "dynamical enhancement" of a weak interaction in the nucleus (see Refs. /31,33,34,30,7/). It is clear that any small interaction in the nucleus is enhanced in a similar way (Ref. /30/). For example, the dynamical enhancement of the Coriolis interaction is considered in the Ref. /35/.

Just as in the case of the weak interaction, the amplitude of the E1-transition between the close compound states turns out to be suppressed /29/. Similarly to /22/, we obtain

$$(E1)_{if} \sim \frac{\Gamma_{sp}}{\sqrt{(\omega-\omega_0)^2 + \Gamma_{sp}^2/4}} \frac{E1\sqrt{q}}{\sqrt{N_i}} \tag{24}$$

Here ω is the frequency of the γ-quantum, E1 is a typical single-particle matrix element.

The estimations we have made above assume a complete incoherence of the contributions of the various terms. This assumption is not always valid. Let us consider, for example, the E1-transition from the ground state of the even-even nucleus to the compound state lying near a giant resonance. In the wave function of this state there are $\sim A^{2/3}$ particle-hole components from which the E1-transition to the ground state is possible (i.e. $q = A^{2/3}$).

In heavy nuclei, the frequency of giant E1-resonance is $\omega \approx 2\omega_0$. Hence, the estimate (24) could give:

$$(E1)_{oc} \sim \frac{A^{1/3}}{\sqrt{N_c}} E1 \frac{\Gamma_{sp}}{\omega_0} \tag{25}$$

However, due to the existence of a giant resonance which plays a role of one of the states φ_α in the wave function (15), there is the coherent contribution of the particle-hole pairs (which "constitute" φ_α) and correct estimate is

$$(E1)_{oc} \sim \frac{A^{1/3}}{\sqrt{N_c}} E1 \tag{26}$$

If we consider now the γ-transition from the ground state to the compound resonance $|c\rangle$, which is distant, in energy, from the giant resonance, then one can see that the resonance contribution decreases

$$(E1)_{oc} \sim A^{1/3} E1 \frac{\langle 1^-_{res}|H_s|c\rangle}{E_c - E_{res}} \sim \frac{A^{1/3}}{\sqrt{N}} E1 \frac{\Gamma_{sp}}{E_c - E_{res}} \tag{27}$$

The hypothesis by Axel and Brink /36/ (see also Ref. /37/) extends, as a matter of fact, the estimate (27) to the transitions between any compound states. This hypothesis is based on the idea that a giant dipole resonance can be excited above any compound state. Indeed, at a frequency of γ-quanta, which is close to the frequency of giant resonance, the incoherent contribution (25) to the E1-amplitude is suppressed and, hence, the hypothesis seems to be quite natural. But the frequencies of transitions in the (n,γ) re-

action do not exceed 8-9 MeV $\approx \omega_0$, that is much less than the frequency of the E1-resonance. In this case, the factor $\Gamma_{spr}/(E_c - E_{res})$ suppresses the resonance contribution and, therefore, the resonance contribution (27) is not larger than the incoherent (24). Moreover, it is not evident that the description with the help of a giant resonance is correct at the frequencies of γ-quanta, which are typical for the single-particle transitions ($\omega \approx \omega_0$).

In Ref. /32/ the idea of the mixing via a giant 0^--resonance was used for estimation of the matrix element of weak interaction between two close compound states. But we have seen that for E1-matrix elements, the resonance mechanism is dominant at the frequency of the giant resonance only due to the fact that the collective frequency, $\omega_{res} \approx 2\omega_0$ is strongly separated from the single-particle frequencies, $\omega \approx \omega_0$. As for the 0^--resonance, if the latter exists, its frequency is likely to be close to the single-particle one: $\omega_{res} \approx \omega_0$. In view of this, the estimations of the matrix element of weak interaction between the compound states which used the "incoherent" (22) and "resonance" /32/ mechanisms coincide with each other. In this statement the words "resonance" and "incoherent" are in inverted commas. The point is that since the resonance frequency is not separated from the single-particle frequencies, the concept of a resonance practically loses sense. The energies of the majority of particle-hole states which belong to the same shell coincide with each other with an accuracy $\delta E \sim \Gamma_{spr}$. Due to the degeneration we can take any combinations of these states as a basic set. As a consequence of this, the assertion that the "resonance" and "incoherent" estimates are the same, becomes entirely clear: the matrix element between physical states is independent of the basis in which we calculate it (i.e. it is not important whether we use particle-hole excitations or their linear combinations, among which there is the 0^--resonance).

The summary of this section consists in that the matrix elements of a one-body operator between two compound states or between the compound and a few-particle states are suppressed as $1/\sqrt{N_{max}}$ compared to the single-particle

matrix element. With the fixed average parameters, the values of these matrix elements are, apparently, distributed according to the Gauss law.

III. Classification of amplitudes in powers of the parameter $1/\sqrt{N}$ is a method of determination of the dominant mechanism of a reaction

In the preceding section the estimates have been obtained for the wave functions of a discrete spectrum. For their applicability to the transitions to a continuous spectrum, let us normalize the neutron wave function at $E > 0$ by the following way

$$\int_{r<R} |\psi_n(r,E)|^2 dV = 1 \qquad (28)$$

R is the size of the nucleus. Besides the matrix elements, the expressions for the amplitudes of reaction include also the energy denominators $(E - E_c + i\Gamma_c/2)$, where E_c, Γ_c are the energy and width of the compound state. These denominators are less by a factor of N than the typical single-particle energy scale $\delta E \sim \Gamma_{spr} \sim \omega_0 \sim$ MeV.

$$\frac{1}{E - E_c + i\Gamma_c/2} = \frac{\mathcal{D}}{E - E_c + i\Gamma_c/2} \frac{\Gamma_{spr}}{\mathcal{D}} \frac{1}{\Gamma_{spr}} \sim N \frac{\mathcal{D}}{E - E_c + i\Gamma_c/2} \frac{1}{\Gamma_{spr}} \qquad (29)$$

Since we want to keep trace of a very large parameter $\sqrt{N}$, we are not concerned now with the other parameters (for example, Γ_{spr}/ω_0). Based upon the above considerations, there is no difficulty to formulate the general rules for classification of amplitudes in powers of small parameter $1/\sqrt{N} \sim 10^{-2}$–10^{-3}. These rules allow to separate the dominant mechanism of a reaction.

1. Each vertex of an interaction (neutron capture, electromagnetic, weak, Coriolis one, etc.) contains the factor $1/\sqrt{N_{max}}$, where $N_{max} = \max(N_c, N_{c'})$; N_c, $N_{c'}$ are the numbers of the principal components of the nuclear wave functions which enter into the vertex. In the ground state N = 1 and in the excited one N ~ MeV/D.

2. Every Green's function of a compound state $|c\rangle$ (at $|E-E_c+i\Gamma_c/2| \lesssim \mathcal{D}$) gives the factor N_c and the resonance energy dependence $\mathcal{D}/(E-E_c+i\Gamma_c/2)$.

It is worth to emphasize that the rules 1) and 2) yield a mean-square estimate of the amplitude. In principle, the fluctuations can occur, which are capable to change the amplitude relation in each particular case. But because of the large value of parameter $\sqrt{N}$, the probability of the fluctuations which violates the $\sqrt{N}$ -hierarchy of amplitudes is very small.

Let us consider, for example, an elastic scattering of a slow neutron ($kR \ll 1$) on a nucleus in the case when only the channels (n,n) and (n, γ) are open. The amplitude of s--wave scattering looks as follows:

$$f(E) = f_0 - \frac{1}{2k}\sum_c \frac{g\,\Gamma_c^{(n)}(E)}{E-E_c+i\Gamma_c/2} \qquad (30)$$

Here $f_0 = -a$, a is the scattering length, $k=\sqrt{2m_nE}$, g is the factor which appears after averaging over the spin projections of the initial nucleus. Note that the separation of the potential term from the resonance one in formula (30) is ambiguous because at a given energy the contribution of far resonances is not distinguishable from the potential amplitude. Therefore, it is worth-while to take into account in formula (30) only several nearest resonances.

If one is interested only in the powers of parameter $\sqrt{N}$, it is simply to write an estimate of the resonance part f_r plotted in Fig. 2. The diagram contains the vertex of neutron capture ($1/\sqrt{N}$), the vertex of neutron emission ($1/\sqrt{N}$), and the Green's function of the compound state ($N\mathcal{D}/(E-E_c+i\Gamma_c/2)$), i.e.

$$f_r \sim \frac{\mathcal{D}}{E-E_c+i\Gamma_c/2} \qquad (31)$$

Thus, in elastic scattering (without parity violation) the resonance contribution, has, just as a potential one, the zero order in powers of $\sqrt{N}$. It follows from the experimental data on the neutron strength function (see, e.g., Ref. /28/) that far from the compound resonances, i.e. at $E-E_c \sim \mathcal{D}/2$,

$$\frac{f_r}{f_0} \sim \frac{1}{ka}\frac{\Gamma_n}{\mathcal{D}} \sim 0.05 \div 0.5 \qquad (32)$$

This ratio is determined by the other parameters (besides $\sqrt{N}$) involved in the consideration of elastic scattering. For example, according to the model considered in the book /28/.

$$\frac{f_r}{f_0} \sim \frac{g\,\Gamma_{sp\,r}}{\pi m_n R a\left[(E-E_a)^2+\Gamma_{sp\,r}^2/4\right]} \qquad (33)$$

Here R is the nucleus radius, E_a is a position of the nearest single-particle s-level. It is seen that the other dimensionless parameters are not so large as $\sqrt{N}$.

The diagrams where the virtual capture and neutron emission are repeated (see Fig. 1) have the same powers of $\sqrt{N}$ as the main diagram. As has been shown in the first section, taking them into account leads, however, only to a some variation of the parameters of compound resonances.

The $1/\sqrt{N}$ -classification in the case of the reaction (n,γ) looks not so trivial. The amplitude is determined by the diagrams given in Fig. 3. The non-resonance part (Fig. 3a) f_0 corresponds to a direct radiative capture of a neutron into the final compound state $|f\rangle$. Since one vertex of the operator of electromagnetic interaction enters into it, this diagram is proportional to $1/\sqrt{N_f}$. The resonance amplitude f_r (Fig. 3b) describes the same process proceeding through the intermediate compound state $|c\rangle$*). This diagram contains the vertex of neutron capture ($1/\sqrt{N_c}$), the Green's function of the compound state ($N_c\,\mathcal{D}/(E-E_c+i\Gamma_c/2)$), and the vertex of photon emission ($1/\sqrt{N_0}$). As a result

$$\frac{f_r}{f_0} \sim \sqrt{N_f}\,\frac{\mathcal{D}}{E-E_c+i\Gamma_c/2} \qquad (34)$$

*) Note that diagram 3c (which contains the factor $\mathcal{D}/(E-E_c+i\Gamma_c/2)\cdot 1/\sqrt{N_f}$) is, in practice, the part of diagram 3b. It corresponds to the separation from $\langle c|H_{em}|f\rangle$ of the contribution of the single-particle components.

Thus, we can see that even at $E-E_c \sim \mathcal{D}$ in the reaction (n, γ) the resonance mechanism dominates in the transitions to the complex final states ($N_f >> 1$), and only for transitions to the ground state and those close to it ($N_f \sim 1$) the direct and resonance mechanisms give a comparable contribution. Since the transitions to the complex states dominate in the total cross section, the known assertion follows that the total cross section of the reaction (n, γ) is determined by the resonance mechanism (see, e.g., Ref. /38/).

Let us now discuss the parity violating effects in neutron reactions. The effects in neutron optics are the simplest ones. The angle of neutron spin rotation around the direction of its motion in the matter and the difference in the total cross sections for the right-hand and left-hand polarized neutrons can be expressed through the parity violating part of the forward elastic scattering amplitude (f_{pv}):

$$\varphi = \frac{2\pi N_0 l}{k} 2Re f_{pv} \qquad (35)$$

$$\Delta\sigma = \sigma_+ - \sigma_- = \frac{4\pi}{k} 2 Im f_{pv} \qquad (36)$$

Here N_0 is the density of the target atoms, l is the path length of the neutron in the sample. The dominant contribution to f_{pv} comes from the diagrams connected with the mixing by a weak interaction of the levels of the compound nucleus with opposite parity (Fig. 4). They contain three vertices $(1/\sqrt{N})^3$ and two Green's functions (N^2), i.e.

$$f_{pv}^{(4)} \sim \sqrt{N} \frac{\mathcal{D}}{E-E_s+i\Gamma_{s1/2}} \frac{\mathcal{D}}{E-E_p+i\Gamma_{p1/2}} \qquad (37)$$

We have taken into account that the s and p are close compound states and, therefore, $N_s \sim N_p = N$. The neutron scattering on the P-odd potential of the nucleus (Fig. 5) and the neutron scattering on the P-odd potential accompanied by a virtual capture into the resonance (Figs. 6 and 7)* also gives the contribution to f_{pv} . The diagram in Fig. 5 has no $\sqrt{N}$, because it is not connected with the compound states. The diagrams in Figs. 6 and 7 has the zero order in $\sqrt{N}$, as well:

$$f_{pv}^{(6a)} \sim 1 \cdot \frac{1}{\sqrt{N}} \cdot N \frac{\mathcal{D}}{E-E_p+i\Gamma_{p/2}} \frac{1}{\sqrt{N}} = \frac{\mathcal{D}}{E-E_p+i\Gamma_{p/2}} \qquad (38)$$

$$f_{pv}^{(6b)} \sim 1 \cdot \frac{1}{\sqrt{N}} N \frac{\mathcal{D}}{E-E_s+i\Gamma_{s/2}} \frac{1}{\sqrt{N}} = \frac{\mathcal{D}}{E-E_s+i\Gamma_{s/2}} \qquad (39)$$

$$f_{pv}^{(7)} \sim \frac{1}{\sqrt{N}} N \frac{\mathcal{D}}{E-E_s+i\Gamma_{s/2}} \frac{1}{\sqrt{N}} \cdot 1 \cdot \frac{1}{\sqrt{N}} \cdot N \frac{\mathcal{D}}{E-E_p+i\Gamma_{p/2}} \frac{1}{\sqrt{N}} = \frac{\mathcal{D}}{E-E_s+i\frac{\Gamma_s}{2}} \cdot \frac{\mathcal{D}}{E-E_p+i\frac{\Gamma_p}{2}} \qquad (40)$$

Thus, the contribution of diagram 4 is enhanced by a factor of $\sqrt{N}$ compared to that of diagrams 5,6 and 7. It is the dynamical enhancement of weak interaction in heavy nuclei that has been mentioned in the second section. The contribution of the diagrams presented in Fig. 5 has been analysed, for example, in Refs. /39-42/. As to the diagrams in Fig. 4, their contribution has been considered in Ref. /18/ (see also Refs. /19-21, 7/), Fig. 6a and 7 - in /43/.

Note that a simple relation follows from the fact that parity violation is associated with the compound resonances which decay mainly through the γ-channel: $\Delta\sigma_{tot} = \Delta\sigma_\gamma$ ($\Delta\sigma = \sigma_+ - \sigma_-$) . The experimental testing of this relation was made in Refs. /13/ and /14/. If the fission channel is open, then $\Delta\sigma_{tot} = \Delta\sigma_\gamma + \Delta\sigma_f$. The amplitudes considered above can be easily expressed through the parameters of compound resonances. For example, the calculation of the dominant amplitude $f_{pv}^{(4)}$ gives (see Appendix B)

$$f_{pv}^{(4)} = \frac{1}{2k} \sum_{sp} \frac{2g\sqrt{\Gamma_s^{(n)}(E)}\, \langle s|H_w|p\rangle\, i\sqrt{\Gamma_{p1/2}^{(n)}(E)}}{(E-E_s+i\Gamma_{s1/2})(E-E_p+i\Gamma_{p1/2})} \qquad (41)$$

Here $\Gamma_{p1/2}^{(n)}$ is the partial width of decay of the p-wave resonance into the state of a neutron with momentum $j = 1/2$, $\langle s|H_w|p\rangle$ is the matrix element of weak interaction between the compound states.

Let us now consider the amplitudes giving rise the parity violating effects in the (n, γ) reaction. Dominant here is the amplitude connected with a weak mixing of the compound

*) Diagram 7 is, in practice, the part of diagram 4. It corresponds to the separation, from $\langle s|H_w|\rangle$ in Eq. (41), of the contribution of the single-particle components.

states formed after the neutron capture (Fig. 8). It contains the vertex of neutron capture ($1/\sqrt{N_1}$), the vertex of weak interaction ($1/\sqrt{N_1}$), the electromagnetic vertex ($1/\sqrt{N_1}$) and two Green's functions of compound states (N_1^2). As a result, this amplitude is enhanced by a factor of $\sqrt{N}$:

$$f_{PV}^{(8)} \sim \sqrt{N_i}\,\frac{\mathcal{D}}{E-E_s+i\Gamma_s/2}\,\frac{\mathcal{D}}{E-E_p+i\Gamma_p/2} \qquad (42)$$

The smaller contribution comes from the amplitude connected with a weak mixing in the final state (Fig. 9):

$$f_{PV}^{(9)} \sim \sqrt{N_f}\,\frac{\mathcal{D}(E\simeq E_i)}{E-E_i+i\Gamma_i/2}\,\frac{\mathcal{D}(E\simeq E_f)}{E-\omega-E_{f'}+i\Gamma_{f'}/2} \qquad (43)$$

Since $\mathcal{D}$ is exponentially energy-dependent, we indicate the energy at which $\mathcal{D}$ is taken.

Besides $f_{PV}^{(8)}$ and $f_{PV}^{(9)}$, there are a number of diagrams not enhanced (e.g., Fig. 10) and even suppressed as $1/\sqrt{N_f}$ (e.g., the contribution of the diagram in Fig. 11 is proportional to $1/\sqrt{N_f}$).

The total number of angular and polarization correlations associated with the interference of various amplitudes is quite large and only a small part of them has been measured. For example, 17 correlations (8 P-even and 9 P-odd) occur in the (n,γ) reaction near the threshold if we take into account the s- and p-wave neutron capture and the dominant EI and MI γ--transitions /ref. 44/:

$$(\vec{n}_n\vec{n}_\gamma),\ \vec{\sigma}[\vec{n}_n\times\vec{n}_\gamma],\ (\vec{n}_n\vec{n}_\gamma)(\vec{\sigma}[\vec{n}_n\times\vec{n}_\gamma]),\ ((\vec{n}_n\vec{n}_\gamma)^2-1/3),$$
$$\lambda(\vec{\sigma}\vec{n}_\gamma),\ \lambda(\vec{\sigma}\vec{n}_n),\ \lambda[(\vec{\sigma}\vec{n}_n)(\vec{n}_n\vec{n}_\gamma)-\tfrac{1}{3}\vec{\sigma}\vec{n}_\gamma],\ \lambda[(\vec{\sigma}\vec{n}_\gamma)(\vec{n}_\gamma\vec{n}_n)-\tfrac{1}{3}\vec{\sigma}\vec{n}_n] \qquad (44)$$
$$(\vec{\sigma}\vec{n}_\gamma),\ (\vec{\sigma}\vec{n}_n),\ [(\vec{\sigma}\vec{n}_\gamma)(\vec{n}_n\vec{n}_\gamma)-\tfrac{1}{3}\vec{\sigma}\vec{n}_n],\ [(\vec{\sigma}\vec{n}_n)(\vec{n}_n\vec{n}_\gamma)-\tfrac{1}{3}\vec{\sigma}\vec{n}_\gamma],$$
$$\lambda,\ \lambda(\vec{n}_n\vec{n}_\gamma),\ \lambda(\vec{\sigma}[\vec{n}_n\times\vec{n}_\gamma]),\ \lambda[(\vec{n}_n\vec{n}_\gamma)^2-\tfrac{1}{3}],\ \lambda(\vec{n}_n\vec{n}_\gamma)(\vec{\sigma}[\vec{n}_n\times\vec{n}_\gamma])$$

$\vec{n}_n$, $\vec{\sigma}$ are the directions of the momentum and spin of the neutron; $\vec{n}_\gamma$, λ are the direction of the momentum and the helicity of the γ-quantum. These correlations have been

calculated in Ref. /44/ with the dominant diagrams taken into account.

The $1/\sqrt{N}$-classification of the amplitudes of the (n,f) reaction proves to be approximately the same as that of the (n,γ) reaction. The main amplitude (Fig. 12) also has the zero order in $\sqrt{N}$. It contains the neutron capture amplitude ($1/\sqrt{N_1}$), the Green's function of a compound nucleus (N_1) and the fission amplitude ($1/\sqrt{N_1}$)*:

$$f^{(12)} \sim \frac{\mathcal{D}}{E-E_s+i\Gamma_s/2} \qquad (45)$$

The dominant admixture amplitude with parity violation (Fig. 13) is enhanced by a factor of $\sqrt{N}$ (see Refs. /30,7/):

$$f^{(13)} \sim \sqrt{N}\,\frac{\mathcal{D}}{E-E_s+i\Gamma_s/2}\,\frac{\mathcal{D}}{E-E_p+i\Gamma_p/2} \qquad (46)$$

There are also non-enhanced amplitudes (Figs. 14 and 15). The diagram in Fig. 15 corresponds to taking into account the weak interaction at the cold stage of fission (Ref. /54/).

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*) The factor $1/\sqrt{N_1}$ in the fission amplitude reflects the magnitude of admixture of the wave function of a cold deformed nucleus in the wave function of a compound state (Refs. /30,7/). This factor leads to a reasonable estimation of the width of the resonance for the above-barrier fission: $\Gamma_f \sim \frac{1}{N}\cdot\frac{1}{\tau} \sim$ 0.1÷1 eV, $\frac{1}{\tau}\sim\frac{v}{R}\sim$ 100 keV, v is the velocity of motion of the fragments at the fission barrier, R the barrier size, τ the lifetime of a cold state.

Appendix A

Non-diagonal mass operator and fluctuations of γ-widths

Let us first recall the way in which the mean square fluctuation $\delta\Gamma$ of γ-widths is estimated. The electromagnetic width of the compound state $|i\rangle$ can be written down as follows:

$$\Gamma_i = \frac{4}{3}\sum_f |M_{if}|^2 \omega_{if}^3 = \sum_f W_{if}\,\omega_{if}^3 \qquad (A1)$$

Here ω_{if} and M_{if} are the frequency and amplitude of the transition to the final state $|f\rangle$, $W_{if} = \frac{4}{3}|M_{if}|^2$. For simplicity, one assumes that angular momentum J and the parity η of the final states $|f\rangle$ are fixed (if several J and η give the contribution to Γ_γ, one needs to sum over them in the following formulae (A2) and (A4)). Let us represent W_{if} in the form $W_{if} = \overline{W_{if}} + \mathcal{E}_{if}$, where the line denotes averaging over the states $|i\rangle$ with the given parity and angular momentum. In the statistical model the quantity W_{if} is distributed according to the Porter-Thomas law (see, for example, Ref. /28/), i.e. $\overline{\mathcal{E}_{if}^2} = 2(\overline{W_{if}})^2$. Then

$$\overline{\Gamma} = \sum_f \overline{W_{if}}\,\omega_{if}^3$$
$$(\delta\Gamma)^2 = \overline{\Gamma^2} - (\overline{\Gamma})^2 = \sum_f \overline{\mathcal{E}_{if}^2}\,\omega_{if}^6 = 2\sum_f (\overline{W_{if}})^2 \omega_{if}^6 \qquad (A2)$$

The imaginary part of the non-diagonal mass operator (8) is of the form (cf. $Im\,\Sigma_{cc}^{\gamma} = -\Gamma_\gamma/2$):

$$Im\,\Sigma_{\kappa i}^{\gamma} = -\frac{2}{3}\sum_f M_{if} M_{f\kappa}\,(E - E_f)^3 \qquad (A3)$$

Since $|\Sigma_{\kappa i}| \lesssim \mathcal{D}$, only the mixing between close compound states is significant, i.e. one can put $E \approx E_\kappa \approx E_i$. After averaging over the states $\kappa, i\ (\kappa \neq i)$ we get $\overline{M_{if} M_{f\kappa}} = 0$,

$$\overline{Im\,\Sigma_{\kappa i}^{\gamma}} = 0$$
$$\overline{\left(Im\,\Sigma_{\kappa i}^{\gamma}\right)^2} = \frac{1}{4}\sum_f (\overline{W_{if}})^2 \omega_{if}^6 = \frac{1}{8}(\delta\Gamma)^2 \qquad (A4)$$

Since the number of states $|f\rangle$ is large, the width fluctuation can be written down as follows:

$$\frac{(\delta\Gamma)^2}{\overline{\Gamma}^2} = \frac{2\sum_J \int_0^{E_i} \{(S_{E1})^2 \rho_{J\eta} + (S_{M1})^2 \rho_{J\bar{\eta}} + \ldots\}\,\omega^6 d\omega}{\left[\sum_J \int_0^{E_i} \{S_{E1}\rho_{J\eta} + S_{M1}\rho_{J\bar{\eta}} + \ldots\}\,\omega^3 d\omega\right]^2} \qquad (A5)$$

Here $S = S(\omega, E_i, J)$ is the radiation strength function, $\rho = \rho(E_i - \omega)$ is the corresponding density of the final states, $\bar{\eta} = -\eta$. For numerical estimation, one can restrict oneself by the dominant multipole (E1,M1) and use the standard parametrization of $\rho(E)$ (see Ref. /28/) with the parameters fitted to the spectrum of a particular nucleus. The frequency dependence of the strength function is weak as compared to the dependence of $\rho(E_i - \omega)$ (for example, $S_{E1} \sim [(\omega-\omega_0)^2 + \Gamma_{sp}^2/4]^{-1}$, according to Eq. (24)) and the former may be neglected. In particular, for ^{117}Sn, we obtain $\delta\Gamma/\Gamma \sim 0.05$.

Appendix B

Calculation of reaction amplitudes

There is no difficulty to write down the calculation rules for reaction amplitudes near the neutron threshold. Let $\vec{n}_\kappa$ be the motion direction of the neutron, $\vec{I}$ the momentum of the initial nucleus, $\vec{J}=\vec{I}+\vec{j}$ the momentum of the compound resonance, α polarization of the neutron, $\vec{j}=\vec{l}+\vec{s}$ the momentum of the p-wave neutron at which the capture occurs. Then,

I. The capture amplitude of the neutron into s-resonance is equal to

$$C^{JJ_z}_{II_z\frac{1}{2}\alpha}\,\eta_s\sqrt{\Gamma_s^{(n)}(E)} \qquad (B1)$$

2. The capture amplitude of the neutron into the p-resonance is equal to

$$\sum_{jj_z m} C^{JJ_z}_{II_z jj_z}\, C^{jj_z}_{1m\frac{1}{2}\alpha}\sqrt{4\pi}\,Y^*_{1m}(\vec{n}_\kappa)\,i\eta_j\sqrt{\Gamma_{p_j}^{(n)}(E)} \qquad (B2)$$

Here $C^{JJ_z}_{II_z\frac{1}{2}\alpha}$ is the Clebsch-Gordan coefficient; $\Gamma_{p_j}^{(n)}$ is the neutron width corresponding to the emittance of a neutron with momentum j ($\Gamma_p^{(n)}=\Gamma_{p_{1/2}}^{(n)}+\Gamma_{p_{3/2}}^{(n)}$); $\eta=\pm1$ is the sign of the amplitude.

3. The matrix element of weak interaction between two compound states is $\langle s|H_w|p\rangle = iW$, where W is real (we use the standard definition of Y_{lm}-functions in accordance with Ref. /46/).

4. The Green's function of a compound nucleus is

$$\frac{1}{E-E_c+i\Gamma_c/2} \qquad (B3)$$

5. The common factor for the scattering amplitudes is equal to $-1/2K$ (K is the neutron moment).

The factor "i" for the p-wave capture is connected with a phase of free motion of the p-wave. We consider the scattering at $kR \ll 1$ and, therefore, the potential scattering phase is zero. As shown in section 1 additional phase factors

$e^{i\varphi}$ ($\varphi\sim\delta\Gamma/2$, $\delta\Gamma$ is the fluctuation of the total width) connected with the diagonalization of the Green's function appear in the rules 1-3. But this factor may be essential only when the width fluctuation is large, for example, if the fission channel is open.

The rules for the fission amplitudes are formulated in Refs. /30, 47,7/. It is also easy to write the amplitudes for radiation (or absorption) of the γ-quantum from the compound state (see, e.g., Refs./48, 44/).

The rules 1-5 were used in Refs. /18, 30,47,7/. Let us test them in the simplest cases. For example, the forward elastic scattering amplitude near the s-resonance is equal to

$$f(0)=-\frac{1}{2K}\,C^{JJ_z}_{II_z\frac{1}{2}\alpha}\sqrt{\Gamma_s^{(n)}(E)}\cdot\frac{1}{E-E_s+i\Gamma_s/2}\cdot C^{JJ_z}_{II_z\frac{1}{2}\alpha}\sqrt{\Gamma_s^{(n)}(E)} \qquad (B4)$$

After summation over J_z and averaging over I_z, one obtains immediately the standard Breit-Wigner formula (30) with $g=\frac{2J+1}{2(2I+1)}$. Similarly, for the p-resonance

$$f(0)=-\frac{1}{2K}\sum_{\substack{jj_z m\\ \tilde{J}\tilde{j}_z\tilde{m}}} C^{JJ_z}_{II_z jj_z}\, C^{jj_z}_{1m\frac{1}{2}\alpha}\sqrt{4\pi}\,Y^*_{1m}(\vec{n}_\kappa)\sqrt{\Gamma_{p_j}^{(n)}(E)}\;\times \qquad (B5)$$

$$\times\frac{1}{E-E_p+i\Gamma_p/2}\cdot C^{JJ_z}_{II_z\tilde{J}\tilde{j}_z}\, C^{\tilde{j}\tilde{j}_z}_{1\tilde{m}\frac{1}{2}\alpha}\sqrt{4\pi}\,Y_{1\tilde{m}}(\vec{n}_\kappa)\sqrt{\Gamma_{p\tilde{j}}^{(n)}(E)}$$

After summation over J_z and averaging over I_z one obtains again the known formula

$$f(0)=-\frac{1}{2K}\,\frac{g\Gamma_p^{(n)}(E)}{E-E_p+i\Gamma_p/2} \qquad (B6)$$

Note that $\Gamma_p^{(n)}=\Gamma_{p_{1/2}}^{(n)}+\Gamma_{p_{3/2}}^{(n)}$.

Finally, the forward elastic scattering amplitude with parity violation is equal to

$$f_{pv}^{(4)}(0)=-\frac{2}{2K}\,C^{JJ_z}_{II_z\frac{1}{2}\alpha}\,\eta_s\sqrt{\Gamma_s^{(n)}(E)}\,\frac{1}{E-E_s+i\Gamma_s/2}\,\langle s|H_w|p\rangle\;\times \qquad (B7)$$

$$\times(-i)\sum_{jj_z m} C^{JJ_z}_{II_z jj_z}\, C^{jj_z}_{1m\frac{1}{2}\alpha}\sqrt{4\pi}\,Y_{1m}(\vec{n}_\kappa)\,\eta_j\sqrt{\Gamma_{p_j}^{(n)}(E)}\,\frac{1}{E-E_p+i\Gamma_p/2}$$

After summation over J_z and averaging over I_z, one obtains

$$f_{PV}^{(A)}(0) = \pm \frac{1}{2K} \frac{2g\sqrt{\Gamma_s^{(n)}(E)}\,(i\langle s|H_w|p\rangle)\sqrt{\Gamma_{p_{1/2}}^{(n)}(E)}\,\eta_s\eta_{1/2}}{(E-E_s+i\Gamma_s/2)(E-E_p+i\Gamma_p/2)} \qquad (B8)$$

The sign (+) corresponds to the positive-helicity neutron, the sign (-) corresponds to the negative-helicity one. After summation the term with $j=3/2$ disappears. The sign factor $\eta_s\eta_{1/2}$ can be excluded by means of the redifinition of the states s and p (i.e. if we introduce it into the matrix element $\langle s|H_w|p\rangle$). Thus we arrive at formula (41).

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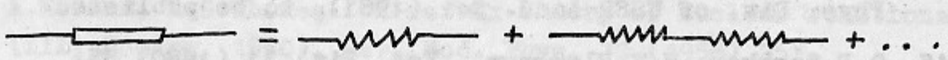


Fig. 1.

— single-particle neutron state

$\sim\sim\sim$ $- (E - E_c + i0)^{-1}$ - Green's function of compound state

$=$ $- (E - E_c - \sum(E))^{-1}$ - "exact" Green's function

Fig. 2.

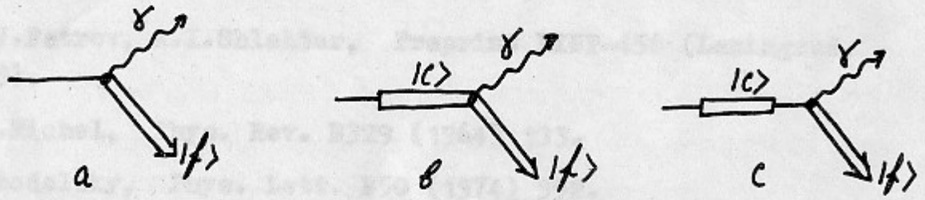


Fig. 3.

Fig. 4.

x - matrix element of weak interaction

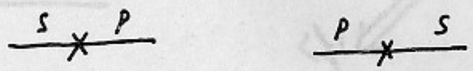


Fig. 5.

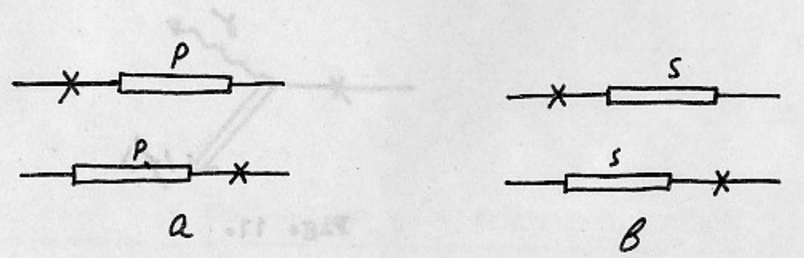


Fig. 6.

Fig. 7.

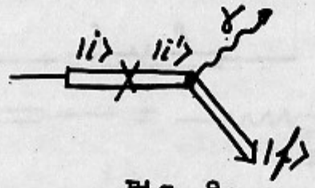


Fig. 8.

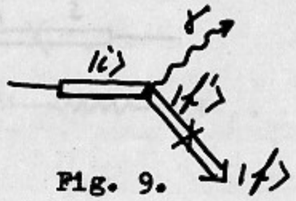
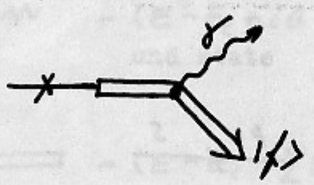


Fig. 9.



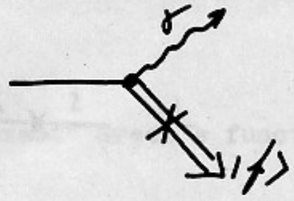


Fig. 10.

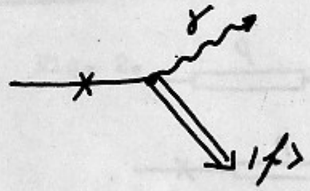


Fig. 11.

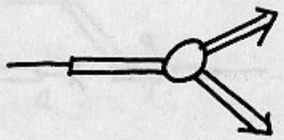


Fig. 12

O - cold stage of fission

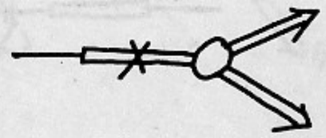


Fig. 13.

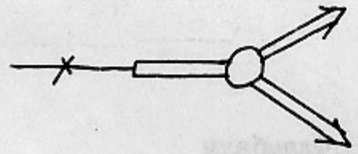


Fig. 14.

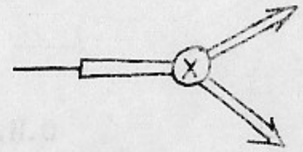


Fig. 15.

О.П.Сушков, В.В.Фламбаум

ТЕОРИЯ НЕЙТРОННЫХ РЕАКЦИЙ И СЛАБОЕ ВЗАИМОДЕЙСТВИЕ В ТЯЖЕЛЫХ ЯДРАХ

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