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FEASIBLE MECHANISM OF PARITY VIOLATION IN NUCLEAR FISSION

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Abstract

The p-odd asymmetry in the angular distribution of fragments from the polarized nucleus fission is considered. A mechanism of this effect is proposed. Asymmetry arises due to mixing of the rotational states of a cold, strongly deformed nucleus. The experiments which enable one to verify this mechanism are discussed.

The spatial parity violation has been observed in the fission of ^{233}U, ^{235}U, ^{239}Pu induced by the polarized thermal neutrons /1-4/. The effect is the correlation of the light fission fragment momentum with the spin of the initial neutron. The asymmetry is $\sim 10^{-4}$.

When attempting to obtain the theoretical treatment of this effect the following problems arise.

1. How does the two-particle weak interaction affect the collective, actually macroscopic motion of the system? Another words, this question may be stated as follows. The fission fragments are heavy. Meanwhile, the mass usually is in the denominator of the effective Hamiltonian of weak interaction. How is this suppression compensated?

2. As known, the number of final states of fission fragments is very large. Why doesn't this circumstance lead to statistical averaging of the effect?

In this study it is proposed a mechanism of the p-odd asymmetry in fission, which gives, apparently, the answer to these questions. The experiments allowing the verification of this mechanism are discussed. It may turn out that the experimental results allow to relate the magnitudes of the effect in different nuclei. As far as the estimation of the effect is concerned, within the framework of the scheme proposed the latter reduced to calculations of certain matrix elements. These calculations appear to be possible but we do not make them here. However, already now an impression is being gained that the theoretical treatment of the p-odd effect in fission is not more complicated than that in the γ-quantum radiation processes.

To illustrate the mechanism of appearing the p-odd correlation in fission, in Appendix I we present a simple model, namely:

the polarized atom ionization. The angular correlation of the decay products due to the mixing of quasi-stationary states is found in Appendix 2.

As known, at not too high energies the nuclear fission process consists of the following stages (see Fig.1):

1. Capture of the particle (for example, the neutron) and formation of a hot compound nucleus.
2. The nucleus is stretched, gets cold and settle down to a certain quantum quasi-stationary state. In other words, the process goes through the fixed fission channel.
3. Breakage of the neck and diverging of the fragments. All the variety of final states are formed just at this stage.

As the experiments show, the p-even angular distributions of the fragments are formed at the stage 2 /5/. This means, first, that the nucleus here lives long enough to "forget" its past and, second, that the breakage occurs so rapidly that information about the state 2 does not lost. Characteristic memory time of the mixing of different states depends only on the distance between the energy levels rather than the intensity of the interaction causing the mixing ($\tau \sim 1/\delta\omega$). Therefore, it is quite reasonable to assume that the p-odd asymmetry also arises at the cold stage. In this case, the second, formulated above problem is solved immediately: asymmetry arises before the breakage and is independent of the final state of fragments.

It is natural to suppose that the mass asymmetry of fragments is formed at the cold stage, i.e. the nucleus in this state represents a pear-shaped top. In the adiabatic approximation the spectrum of such a system looks as follows /5/. At a given internal state $|a, K\rangle$ (K is the projection of the total angular momentum J on the top axis) there is a band of rotational states (see

Fig.2). If $K \neq 0$, then at each J there exist two close levels of opposite parity. This is the phenomenon similar to the Λ -doubling in moleculas[1)].

$$|a,K\rangle^{\eta}_{JM} = \sqrt{\frac{2J+1}{8\pi}}\left\{\mathcal{D}^{J}_{MK}(\varphi,\theta,0)\,|a,K\rangle + \eta(-1)^{J+K}\mathcal{D}^{J}_{M,-K}\,|a,-K\rangle\right\} \qquad (1)$$

Here η is the parity of the state. It is easy to verify that in the adiabatic approximation the weak interaction does not mix these levels. Actually, the matrix element ${}^{\eta}_{JM}\langle a,K|H_w|a,K\rangle^{\bar\eta}_{JM}$ ($\bar\eta \equiv -\eta$) reduces to the average value H_w in the body-fixed frame $\langle a,K|H_w|a,K\rangle$. H_w is the pseudoscalar, therefore, $\langle a,K|H_w|a,K\rangle \sim K$. But since K changes its sign at the time reflection, this relation contradicts to the T-invariance of H_w. Hence, ${}^{\eta}_{JM}\langle a,K|H_w|a,K\rangle^{\bar\eta}_{JM} = 0$. The mixing which is not equal to zero is the p-odd mixing with the levels $|b,K\rangle_J$ whose internal states differ from $|a,K\rangle_J$:

$$\widetilde{|a,K\rangle}^{\eta}_{JM} = |a,K\rangle^{\eta}_{JM} + \sum_{b}\frac{\langle b,K|H_w|a,K\rangle}{E_a - E_b}\,|b,K\rangle^{\bar\eta}_{JM} \qquad (2)$$

However, this mixing itself does not yet lead to the angular asymmetry since the wave functions $|a\rangle$ and $|b\rangle$ do not interfere due to orthogonality over internal variables. In order for the interference to appear it is necessary to return the system from the state $|b\rangle$ to $|a\rangle$. This transition can be caused by any interaction H_K violating adiabaticity. There is no difficulty to see that for the effect under discussion only the part of the operator H_K which changes the sign K is significant, i.e. $|b,K\rangle \to |a,-K\rangle$. Thus, if one takes into account the p-odd mixing, the wave function of the state $|a\rangle$ has the form:

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1) Such a rotational structure has been observed in neutron-induced fission of ^{230}Th /6/.

$$|\tilde{a,K}\rangle^{\eta}_{\mathcal{J}M} = |a,K\rangle^{\eta}_{\mathcal{J}M} + \frac{i\eta u_w}{E_\eta - E_{\bar\eta}}\,|a,K\rangle^{\bar\eta}_{\mathcal{J}M}$$

$$i\,u_w = -2\,A(\mathcal{J},K)\sum_{\ell}\frac{\langle a,-K|H_K|\ell,K\rangle\langle\ell,K|H_w|a,K\rangle}{E_a - E_\ell} \qquad (3)$$

u_w is real. The coefficient $A(\mathcal{J},K)$ depends on the transformation properties of H_K:

$${}^{\eta}_{\mathcal{J}M}\langle a,K|\,H_K\,|\ell,K\rangle^{\eta}_{\mathcal{J}M} = \eta\,A(\mathcal{J},K)\langle a,-K|H_K|\ell,K\rangle$$

Energy splitting of the states η and $\bar\eta$ is also due to the interaction H_K:

$$E_\eta - E_{\bar\eta} = 2\eta\,A_{\mathcal{J}K}\,\langle a,-K|\,H_K|a,K\rangle$$

We do not know a concrete form of the operator H_K. The part of H_K can take the Coriolis interaction, the interaction of internal degrees of freedom with flexural oscillations of the stretched nucleus (apparently, such oscillations occur in fission; see, for example, Ref./7/) and finally, the interaction corresponding to the places exchange of the light and heavy fragments /5/ (in this case the transition can be above-barrier). Some information on H_K can be extracted from the measurements of energy splitting $E_\eta - E_{\bar\eta}$.

For rough estimation of the mixing coefficient $\beta = i\eta u_w/(E_\eta - E_{\bar\eta})$ the matrix elements of H_K in both the numerator and denominator can be cancelled. If one assumes that $|a\rangle$ and $|\ell\rangle$ are the single-particle states and also that $E_a - E_\ell \sim 1$ MeV, then, as usual, $\beta \sim 10^{-6}$. This is considerably smaller than the observed value of the p-odd asymmetry. Our hopes for amplifying the effect are associated with that, first, the distance between the levels in strongly deformed nuclei is much smaller than that in the non-deformed ones. Second, amplification is possible due to collective

character of the matrix element $\langle\ell|H_w|a\rangle$. The point is that the state $|a\rangle$ is not the single-particle one. Apparently, it is a collective state with $K = 1$ which is analogous to a low-lying branch of the octupole excitation in symmetrically deformed nuclei[2].

Besides the mixing coefficient, asymmetry also depends on the decay amplitudes of the states $|a\rangle^{\eta}$ and $|a\rangle^{\bar\eta}$. For example, if $|a\rangle^{\eta}$ and $|a\rangle^{\bar\eta}$ are the bound states, then the mixing coefficient is purely imaginary due to T-invariance of the operator H_w. For this reason the interference between these states does not appear. It is quite natural and corresponds to the absence of electrical dipole moment in the stationary state. Interference arises in fission due to that the decay phases of $|a\rangle^{\eta}$ and $|a\rangle^{\bar\eta}$ are different. The angular correlation of the decay products due to the mixing of quasi-stationary states is found in Appendix 2. The angular distribution of fragments in fission through the channel $|a,K\rangle$ from the state with a given energy E, angular momentum $\mathcal{J}$, $\mathcal{J}_z = M$, and parity η has the form:

$$W_{\mathcal{J}M}(\theta) \sim |\mathcal{D}^{\mathcal{J}}_{MK}|^2(1+\gamma) + |\mathcal{D}^{\mathcal{J}}_{M,-K}|^2(1-\gamma) \qquad (4)$$

$$\gamma = 2\eta\,\mathrm{Re}\,\frac{i\,u_w}{E - E_{\bar\eta} + i\Gamma/2}$$

In the experiments the fission of the unpolarized nuclei has been induced by polarized neutrons /1-4/. In this case, the angular distribution is the following:

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2) There is experimental evidence for that the fission of ^{236}U and ^{234}U goes from the state with $K = 1$ /8/. It is very important that $K \neq 0$ since at $K = 0$ the discussed by us mechanism of appearing the asymmetry does not "work".

$$W(\theta) \sim \sum_M |C^{JM}_{IM-1/2,\,1/2\,1/2}|^2 W_{JM}(\theta) \sim 1 + a\cos\theta$$

$$a = \frac{\eta\, U_W \Gamma}{(E-E_{\bar\eta})^2 + \Gamma^2/4}\, \frac{K}{I+1/2} (-1)^{J-I-1/2} \tag{5}$$

I is the angular momentum of the target nucleus, $C^{JM}_{IM-1/2,1/2\,1/2}$ is the Clebsch-Gordan coefficient. If the fission goes through some channels with different K, then $a = \sum_k W_k a_k$, where W_k is the corresponding fission probability.

In our discussions the p-odd correlation is assumed to be caused by the usual weak interaction. In principle, there is another possibility. To the same effect leads the interaction H_{TP} violating both the T- and P-parities simultaneously (namely this interaction is responsible, for example, for a hypothetic electrical dipole moment of a neutron). In this case, formulae (4) and (5) are also applicable in fact. The difference consists in that the imaginary matrix element $i U_W$ should be replaced by the real U_{TP}:

$$a = 2\eta\, \mathrm{Re}\, \frac{U_{TP}}{E - E_{\bar\eta} + i\Gamma/2}\, \frac{K}{I+1/2} (-1)^{J-I-1/2} =$$

$$= 2\eta\, \frac{U_{TP}\,(E-E_{\bar\eta})}{(E-E_{\bar\eta})^2 + \Gamma^2/4}\, \frac{K}{I+1/2} (-1)^{J-I-1/2} \tag{6}$$

Unlike the matrix element of the weak interaction (3), U_{TP} is not equal to zero even with no taken into account the interaction H_K which breaks adiabaticity, i.e. already in the first order of perturbation theory $U_{TP} = \langle a, K | H_{TP} | a, K \rangle$. This circumstance can partially compensate a decrease in the effect value which is connected with the smallness of the interaction H_{TP} itself. Though, it is extremely hardly probable that the observed effect is caused by this interaction. In principle, the question of what interaction causes the angular asymmetry may be solved experimentally since the energy dependence of the effect in formulae (5) and (6) is different.

Return now to the questions which have been posed at the beginning of this paper. As we see, the answer to the first one is that the smallness of Γ and the interval between the rotational levels of opposite parity compensate the suppression of the effect due to a large nucleus mass. The answer to the second question has already been formulated. It consists in that the asymmetry is formed at the cold stage of fission.*)

The effect magnitude is not estimated in the present paper. But the scheme proposed makes it possible to point out the experiments which, on the one hand, allow one to verify it and, on the other hand, to determine the parameters necessary for qualitative predictions.

1. It is of significance to determine J and K of the fission channel where the effect is observed. As known, to this end the usual p-even correlations in the angular distribution of fragments should be measured. But for the corresponding experiment a polarized target is required (as mentioned above, for ^{235}U and ^{233}U some experimental information is already available). If it turns out that in the nuclei where the effect was observed, the same fission channel works (This seems to be quite reasonable), the possibility will arise to relate the magnitudes of p-odd asymmetry for these nuclei.

*) It is very important that both the ground and mixed levels in (3) have the same internal state. Therefore, the ratio of fission amplitudes from the levels of opposite parity doesn't depend on a concrete internal state of fragments.

2. It is of interest to measure the dependence both of p-odd and usual p-even angular distributions on the energy of the neutron. Such information may allow to determine $\mathcal{U}_w$, Γ and $E_+ - E_-$.

3. In the discussed model the p-odd asymmetry is independent of the nuclear excitation method. Therefore, in fission induced by any other particles, for example, in photofission of ^{234}U, ^{236}U, ^{240}Pu, the effect arises too. Its value can be predicted if the fission goes through the same channel as that induced by neutrons.

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APPENDIX 1

POLARIZED ATOM IONIZATION

Ionization of a polarized atom is the simplest model of arising the p-odd asymmetry in the angular distribution of fission fragments. Let us consider a hydrogen atom in the ground state (for the sake of simplicity the nucleus is assumed to be spinless). Due to the weak interaction, the p-waves are admixed to the state $1S$:

$$\Psi_i = |S_{1/2}, J_z = \tfrac{1}{2}\rangle + \frac{\langle P_{1/2}|H_w|S_{1/2}\rangle}{E_S - E_P} |P_{1/2}, J_z = \tfrac{1}{2}\rangle \qquad \text{A1.1}$$

It is easy to see that the mixing coefficient is purely imaginary: $\langle P_{1/2}|H_w|S_{1/2}\rangle/(E_s-E_p) = i\alpha$. Let at a ceratin moment the nucleus disappear instantly, i.e. the electron becomes free:

$$\Psi_f = e^{i\vec{k}\vec{r}}\chi = 4\pi \sum_{lm} i^l j_l(kr) Y^*_{lm}\left(\frac{\vec{k}}{k}\right) Y_{lm}\left(\frac{\vec{r}}{r}\right)\chi , \qquad \text{A1.2}$$

χ is the spin function. The transition amplitude $i \to f$ has the following form:

$$A = \langle \Psi_f|\Psi_i\rangle \sim \langle j_0 Y_{00}\chi|S_{1/2}\rangle Y_{00}\left(\frac{\vec{k}}{k}\right) + \alpha \langle j_1 Y_{1m}\chi|P_{1/2}\rangle Y_{1m}\left(\frac{\vec{k}}{k}\right) \qquad \text{A1.3}$$

Turn out attention to the fact that the imaginary unit in the mixing coefficient in formula A1.1 is compensated by the imaginary unit in the expansion of the exponent in formula A1.2. Since the factor i^l in formula A1.2 is connected with the free motion phases of spherical waves, we can say that the imagineryness of the mixing coefficient is compensated due to difference in the motion phases of S- and P-states.

The angular distribution of the outgoing electrons looks as follows:

$$W(\theta) \sim 1 - \alpha \frac{A_P}{A_S} \cos\theta, \qquad \text{A1.4}$$

A_P and A_S are the overlapping of the corresponding radial wave functions of the free and bound electrons.

In order to avoid misunderstanding, it should emphasize that the model under study highly differs from the real fission process since ionization here occurs under the instant external perturbation. For this reason, formula A1.4 differs from formula A2.1.

APPENDIX 2

INFLUENCE OF QUASI-STATIONARY STATES MIXING ON THE ANGULAR DISTRIBUTION OF DECAY PRODUCTS

Nuclear fission starts from a certain quasi-stationary state. Therefore, from the standpoint of quantum mechanics we are interested in the following problem: a particle is produced inside the potential well, near the quasi-stationary level, and left it some time after. A graph of this process is shown in Fig.3a. The cross denotes the matrix element H_S describing a capture of the neutron and also a transition of the system through the compound-nucleus stage into the cold state. The admixture amplitude is plotted in Fig.3b. The point denotes the matrix element of the p-odd mixing $i\eta U_W$ (3). The decay amplitude of the state $|\eta\rangle$ can be written in the form $A=\sqrt{\Gamma_\eta}\exp(i\varphi_\eta)$. Then the wave function of the final state is the following (Γ occurs in the denominator after summation of diagrams describing the virtual decays):

$$\Phi_E \sim \frac{H_S}{E-E_\eta+i\Gamma_\eta/2}\sqrt{\Gamma_\eta}\,e^{i\varphi_\eta}\Theta_\eta + \frac{H_S}{(E-E_\eta+i\Gamma_\eta/2)}\cdot\frac{i\eta U_W}{(E-E_{\bar\eta}+i\Gamma_{\bar\eta}/2)}\sqrt{\Gamma_{\bar\eta}}\,e^{i\varphi_{\bar\eta}}\Theta_{\bar\eta} \sim$$
$$\sim \Theta_\eta + \frac{i\eta U_W}{E-E_{\bar\eta}+i\Gamma_{\bar\eta}/2}\sqrt{\frac{\Gamma_{\bar\eta}}{\Gamma_\eta}}\,e^{i(\varphi_{\bar\eta}-\varphi_\eta)}\Theta_{\bar\eta} \qquad \text{A2.1}$$

Here Θ_η and $\Theta_{\bar\eta}$ are the angular wave functions corresponding to the final states of opposite parity. For the fission model under discussion the states $|\eta\rangle$ and $|\bar\eta\rangle$ correspond to the same internal state of the nucleus. Therefore, $\Gamma_\eta=\Gamma_{\bar\eta}=\Gamma$, $\varphi_{\bar\eta}=\varphi_\eta$ so that

$$\Phi_E \sim \Theta_\eta + \frac{i\eta U_W}{E-E_{\bar\eta}+i\Gamma_{\bar\eta}/2}\,\Theta_{\bar\eta} \qquad \text{A2.2}$$

The function Θ_η will be given below (see A2.15).

In order to clarify a picture of the process and physical approximations, it seems to be reasonable to derive the formula A2.2 another, though much more cumbersome way, namely: we shall find a wave function describing the decay of the quasi-stationary state.

In our problem the boundary condition at $r \to \infty$ is a divergent spherical wave. However, it is convenient to construct a perturbation theory, basing upon the standing spherical waves (diverging waves are singular in zero). Remind that in this basis a diverging wave occurs in consideration of the wave packet.

$$\psi(r,t) = \frac{1}{r}\int_0^\infty \psi_k \sin(kr+\delta_k)\, e^{-i\frac{k^2}{2m}t}\, dk \qquad \text{A2.3}$$

At $r, t \to \infty$ the integral is calculated according to the stationary phase method. There is no difficulty to see that in this case in the sinus only the diverging wave is significant. Thus, the recipe consists in the following. All calculations can be made for standing waves. In order to come to the real problem in the wave functions obtained, at $r \to \infty$ the replacement should be performed:

$$\sin(kr+\delta_k) \to \frac{1}{2i} e^{i(kr+\delta_k)}$$

Let us now consider the perturbation theory for quasi-stationary states. To this end, we shall consider a particle motion with some angular momentum. Let there exist a number of quasi-stationary levels with energies E_n and widths Γ_n in this wave. The radial wave function of the state of continuous spectrum with energy E may be written in the following form:

$$\psi_E(r) = \sin(kr+\delta_E) \quad \text{at} \quad r \to \infty \qquad \text{A2.4}$$

$$\psi_E(r) = A_E\, \chi_E(r) \quad \text{at} \quad r \lesssim r_0 \qquad \text{A2.5}$$

r_0 is the well size. The wave function $\chi_E(r)$ is real and normalized

$$\int_{r \lesssim r_0} \chi_E^2(r)\, dr = 1,$$

i.e. at $E = E_n$, $\Gamma_n \to 0$ χ becomes a wave function of discrete spectrum. Near the resonance

$$A_E = \frac{1}{\mathcal{D}_n} \frac{\Gamma_n/2}{\sqrt{(E-E_n)^2 + \Gamma_n^2/4}}, \qquad \text{A2.6}$$

$$\delta_E = -\frac{\pi}{2} - \operatorname{arctg}\frac{2(E-E_n)}{\Gamma_n} + \varphi_E, \qquad \text{A2.7}$$

φ_E is the non-resonant part of the phase. The coefficient $\mathcal{D}_n$ is defined from the normalization condition

$$\int A_E^2 \frac{2}{\pi} dk = 1 \qquad \text{A2.8}$$

and is $\mathcal{D}_n = \sqrt{\Gamma_n m / k_n}$. The element of the phase volume in Eq. A2.8 is determined by normalization of the function A2.4. Let us now include the perturbating potential $U(r)$. It is convenient to write Schredinger's equation for the perturbed function in the integral form:

$$\tilde{\psi}_E = \psi_E + \int_0^\infty \frac{\langle \psi_{E'} | U | \tilde{\psi}_E \rangle}{E - E'} \psi_{E'} \frac{2}{\pi} dk' \qquad \text{A2.9}$$

Since we want to deal with the basis of standing waves, by the integral in this formula the principal value is meant[3].

3) An integral is usually taken with the denominator $E - E' + i\varepsilon$ that corresponds to appearing a diverging wave at infinity. We prefer not to leave the basis of standing waves. Of course, when passing to the wave packets, both methods give the same answer.

Let the energy E be close to that of one of the resonances, for example, the first. It is necessary for us to know the wave function $\tilde{\psi}_E$ at $r \le r_0$ and at $r \to \infty$. In both cases, the integral in Eq.A2.9 is calculated explicitly. Actually, the main contribution to the matrix element $\langle \psi_{E'} | U | \tilde{\psi}_E \rangle$ is made by the region wherein the wave functions are amplified by a resonance. Therefore, the dependence of $\langle \psi_{E'} | U | \tilde{\psi}_E \rangle$ on E' is determined by the resonant factor $A_{E'}$. At $r \le r_0$, when the wave function $\psi_{E'}$ has the form of Eq.A2.5, the integral is divided into the sums of resonant contributions. At $r \to \infty$ the function $\psi_{E'}$ oscillates quickly and the integral is determined by the pole contribution at $E' \simeq E$. Hence, the integral equation A2.9 reduces to algebraic equations. At $\Gamma, E-E_1, U \ll E-E_{n\neq 1}$ the solution is the following:

$$\tilde{\psi}_E = \sqrt{Z}\sin(kr+\delta_E) \quad \text{at} \quad r \to \infty$$

$$\tilde{\psi}_E = \sqrt{Z}\tilde{A}_E \tilde{\chi}_E \quad \text{at} \quad r \le r_0$$

$$\tilde{\chi}_E = \chi_{E-\langle \chi_1|U|\chi_1\rangle} + \sum_{n\neq 1} \frac{\langle \chi_n|U|\chi_1\rangle}{E-E_n}\chi_n \qquad \text{A2.10}$$

$$\tilde{\delta}_E = -\frac{\pi}{2} - \mathrm{arctg}\,\frac{2(E-\tilde{E}_1)}{\Gamma_1} + \varphi_E$$

$$\tilde{A}_E = \frac{1}{D_1}\,\frac{\Gamma_1/2}{\sqrt{(\tilde{E}-E_1)^2 + \Gamma_1^2/4}}$$

$$\tilde{E}_1 = E_1 + \langle \chi_1|U|\chi_1\rangle$$

The factor $\sqrt{Z}$ is not given because it should be eliminated by renormalization of the wave function $\tilde{\psi}_E$ [4]. Formulae A2.10

4) The integral equation A2.9, as well as perturbation theories in the continuous spectrum yields a non-normalized wave function.

may also be obtained by a direct summation of a series of perturbation theories. Just as one should expect, after normalization $\tilde{\psi}_E$ takes the form A2.4-A2.7 and $\tilde{\chi}_E$ at $E=E'$ is expressed in terms of perturbation theory for discrete spectrum. Emphasize that the non-resonant part of the phase in ψ_E and $\tilde{\psi}_E$ is the same.

Return now to nuclear fission. In the adiabatic approximation the opposite-parity states in the nucleus are degenerated and the wave functions have the form (1). The interaction H_K breaking adiabaticity causes the energy splitting and results in mixing with other internal states. As shown above (see A2.10), at $r \lesssim r_0$ the wave function of continuous spectrum looks as follows:

$$\psi_E^\eta = A_E^\eta \chi_\eta$$

$$\chi^\eta = |a,K\rangle^\eta_{JM} + \sum_b \frac{{}_{JM}\langle b,K|H_K|a,K\rangle^\eta_{JM}}{E_a - E_b}\,|b,K\rangle^\eta_{JM} \qquad \text{A2.11}$$

The weak interaction leads to mixing of the states ψ^η and $\psi^{\bar{\eta}}$:

$$\langle \psi_E^{\bar{\eta}}|H_w|\psi_E^\eta\rangle = i\eta u_w A_E^{\bar{\eta}} A_E^\eta \qquad \text{A2.12}$$

The matrix element $\langle \chi^{\bar{\eta}}|H_w|\chi^\eta\rangle \equiv i\eta u_w$ is given in Eq.(3). The integral in

$$\tilde{\psi}_E^\eta = \psi_E^\eta + \int \frac{\langle \psi_{E'}^{\bar{\eta}}|H_w|\psi_E^\eta\rangle}{E-E'}\,\psi_{E'}^{\bar{\eta}}\,\frac{2}{\pi}\,dk' \qquad \text{A2.13}$$

is calculated in the same way as that in Eq.A2.9. Taking into account the p-odd mixing, we get the following wave functions:

$$\tilde{\psi}_\eta = A_E^\eta \left\{ \chi^\eta + \frac{i\eta u_w \Delta_{\bar{\eta}}}{\Delta_{\bar{\eta}}^2 + \Gamma^2/4}\,\chi^{\bar{\eta}} \right\} \quad \text{at} \quad r \lesssim r_0 \qquad \text{A2.14}$$

$$\tilde{\Psi}_E^\eta = \Theta^\eta \cdot \sin(kr + \delta_E^\eta) - \frac{i\eta\, U_w\, \Gamma/2}{\sqrt{\Delta_\eta^2 + \Gamma^2/4}\sqrt{\Delta_{\bar\eta}^2 + \Gamma^2/4}} \cos(kr + \delta_E^{\bar\eta}) \cdot \Theta^{\bar\eta} \quad at\ r \to \infty$$

$$\delta_E^\eta = -\frac{\pi}{2} - \mathrm{arctg}\,\frac{2\Delta_\eta}{\Gamma} + \varphi, \quad \Delta_\eta = E - E_\eta$$

At $r \to \infty$ in the wave functions we single out the factors responsible for a relative motion of fragments:

$$|a,K\rangle^\eta_{JM} \to \Theta^\eta \cdot \sin(kr + \delta^\eta) \quad \text{at} \quad r \to \infty$$

$$\Theta^\eta = \sqrt{\frac{2J+1}{8\pi}} \left\{ \mathcal{D}^J_{MK} |\tilde{a},K\rangle + \eta \cdot (-1)^{J+K} \mathcal{D}^J_{M,-K} |\tilde{a},-K\rangle \right\} \qquad \text{A2.15}$$

$|\tilde{a},K\rangle$ is the internal wave function for fragments.

In addition to the wave functions $\tilde{\Psi}^\eta$ describing both the cold stage of fission and divergence, it is necessary to know the transition amplitudes of the system to these states. Let a certain operator H_S describe a capture of the neutron and transition of the system via compound-nucleus state into the cold state. As has been said above, the weak interaction in the compound nucleus is assumed to be negligible, i.e. H_S conserves parity. If one takes into account the fact that parity of the initial state Ψ_i is fixed, one gets the wave function of motion of the fragments:

$$\Phi(\vec{r}) = \sum_\alpha \langle \tilde{\Psi}^\alpha | H_S | \Psi_i^\eta \rangle \tilde{\Psi}^\alpha = \langle \Psi^\eta | H_S | \Psi^\kappa \rangle \sum_\alpha \langle \tilde{\Psi}^\alpha | \Psi^\eta \rangle \tilde{\Psi}^\alpha \sim$$

$$e^{i(kr+\delta^\eta)} \left\{ \Theta^\eta + \frac{i\eta\, U_w (\Delta_\eta - i\Gamma/2)}{\sqrt{\Delta_\eta^2 + \Gamma^2/4}\sqrt{\Delta_{\bar\eta}^2 + \Gamma^2/4}} e^{i(\delta^{\bar\eta} - \delta^\eta)} \Theta^{\bar\eta} \right\} \qquad \text{A2.16}$$

We here remain only the terms corresponding to a diverging wave. Substituting the phase difference from A2.14, we obtain the wave

angular function A2.2.

Turn our attention to an interesting circumstance: the p-odd correlation in A2.2 is proportional to

$$\mathrm{Re}\, \frac{i U_w}{E - E_{\bar\eta} + i\Gamma/2} \sim \Gamma$$

There is no difficulty to clarify this fact, basing on T-invariance of the weak interaction. The matter is that the state of A2.2 and A2.16 with the asymmetric diverging wave has a non-zero dipole moment directed along the total moment of the system $\vec{J}$. This dipole moment is, of course, proportional to a width. Thus, we see a concrete example which shows that the widely discussed dipole moment of the unstable state /9/ (see also Refs./10-12/) is none other than a usual p-odd correlation in the angular distribution of decay products.

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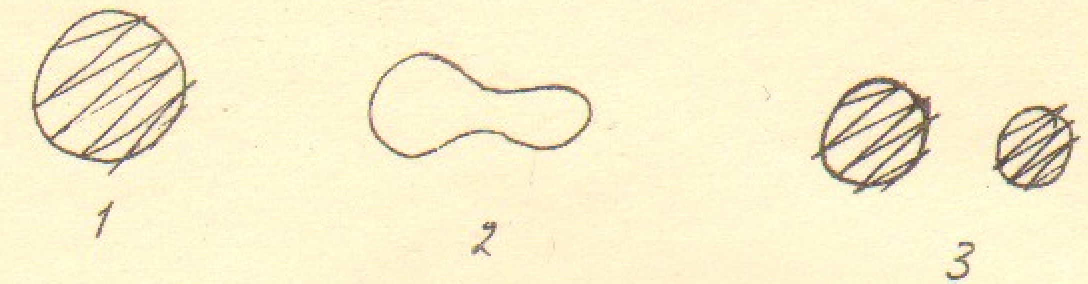

Fig.1

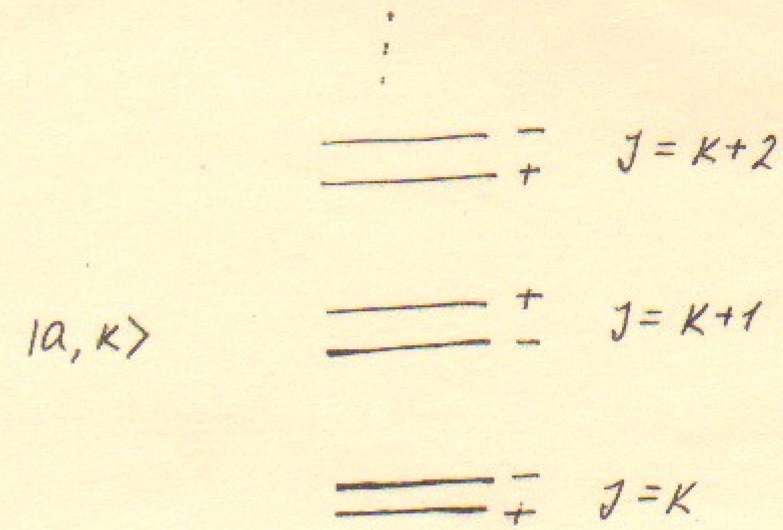

Fig.2

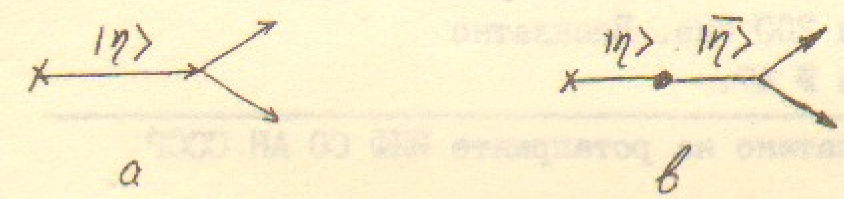

Fig.3