# Parity breaking in the interaction of neutrons with heavy nuclei

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Mechanisms for the following parity-breaking effects are examined: a) asymmetry in the emission direction of the light fission fragment; b) rotation of the neutron spin around the direction of its momentum in a medium and the difference between the cross sections for the capture of right-hand- and left-hand-polarized neutrons; c) circular polarization and asymmetry in the angular distribution of $\gamma$ rays. There is particular emphasis on the reasons for the enhancement of these effects in heavy nuclei. Theoretical results are compared with experimental data.

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## CONTENTS



## INTRODUCTION

One of the most interesting recent discoveries in nuclear physics has been the observation of parity breaking in the fission of nuclei by polarized neutrons.[1-4] It has been found that the probabilities for the emission of the light fragment are different in the directions along and opposite the spin direction of the initial neutron. The magnitude of this asymmetry is $\sim 10^{-4}$. This discovery was so surprising that the original experiment remained under doubt until it was confirmed.[5-12] The significance of the effect in fission is that, in contrast with the *P-odd* effects which had been observed previously, the parity breaking in the fission case is manifested in the motion of a nearly macroscopic object: a fragment consisting of $10^2$ nucleons. Section 2 of this paper will deal with the theory of this effect. It turns out that the effect arises because the fission process occurs in an extremely unusual manner: Although many degrees of freedom are excited at both the beginning and the end of the fission process, it goes through only a small number of intermediate collective states (fission channels). The *P*-odd angular correlations, like the ordinary *P*-even correlations, arise in this "cold" intermediate step of the fission. In this step the nucleus is a pear-shaped top (in other words, the highly deformed nucleus consists of two clusters with different masses, which are approximately equal to the masses of nuclei with filled shells). The asymmetry in the fragment emission direction results from a mixing of rotational levels of opposite parity in this system. In turn, this mixing results from a weak interaction during the compound-nucleus step. The transfer of the *P*-odd mixing from one step to another is related to the slight uncertainty in the energy of the excited nucleus; the time dependence can thus be separated out from the total wave function, so that the "forgetting" of the initial step can be prevented. The asymmetry predicted on the basis of this mechanism, with the dynamic enhancement in the compound nucleus taken into account, agrees with experiment.

The *P*-even correlation between the momentum of the fission fragment and the normal to the plane defined by the spin and momentum of the incident neutron was measured in some recent experiments.[10,11] It will be shown in Sec. 3 of the present review that the mechanism responsible for this correlation has much in common with that for the occurrence of the *P*-odd asymmetry. The estimated magnitude of this correlation ($\sim 10^{-4}$) agrees with experiment.

Yet another effect in which parity breaking is of a coherent macroscopic nature is seen in neutron optics. It will be shown in Sec. 4 that the corresponding effects are significantly enhanced near *p*-wave compound resonances. The relative difference between the absorption cross sections for right-hand- and left-hand-polarized neutrons may be more than $10^{-2}$. As transverse-

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polarized neutrons pass through a medium, their spin is rotated around their momentum direction. The rotation angle over a mean free path is $10^{-2}$–$10^{-3}$ rad near a $p$-wave resonance and $10^{-4}$–$10^{-5}$ rad for thermal neutrons. As unpolarized neutrons pass through a medium, they acquire a longitudinal polarization; the degree of polarization over a mean free path is $10^{-2}$–$10^{-3}$ near a $p$-wave resonance. Rotation of the spin of thermal neutrons in $^{117}$Sn has already been observed experimentally.[13] The results agree with the theoretical predictions.

Also enhanced near $p$-wave resonances are parity-breaking effects in the $(n,\gamma)$ reaction. In this case the effects are a circular polarization of the $\gamma$ rays [$(\mathbf{s}_\gamma\mathbf{p}_\gamma)$ $(\mathbf{p}_\gamma\mathbf{p}_n)$ and $(\mathbf{s}_\gamma\mathbf{p}_\gamma)$ correlations] and an asymmetry in the $\gamma$ angular distribution. These effects will be discussed in Sec. 5.

We should point out that at present it would hardly be possible to extract any accurate quantitative information about the weak interaction between neutrons from a study of these parity-breaking effects in heavy nuclei, because of the complexity of these systems. In such cases the weak interaction is not so much itself the object of the study as a tool for studying various physical effects in a nucleus: the fission process, dynamic enhancement, etc. For example, the experiments in neutron optics which we will be discussing might yield some extremely comprehensive information about the structure of compound states.

## 1. MAGNITUDE OF THE PARITY BREAKING IN NUCLEAR FORCES. DYNAMIC ENHANCEMENT

The parity-breaking weak interaction of a nucleon in a nucleus can be described approximately by the effective Hamiltonian

$$H_W \sim G\frac{\sigma\mathbf{p}}{2m}\rho, \tag{1.1}$$

where $G=10^{-5}/m^2$ is the Fermi constant; $\sigma$, $\mathbf{p}$, and $m$ are the spin, momentum, and mass of the nucleon; and $\rho$ is the nuclear density. A measure of the mixing of single-particle levels of opposite parity by the weak interaction is the ratio $H_W/\omega$, where $\omega\sim p^2/2m$ is the characteristic energy of the nucleon. In a nucleus we have $p\sim m_\pi$ and $\rho\sim 1/m_\pi^3$, so that

$$F\sim\frac{H_W}{\omega}\sim Gm_\pi^2=2\cdot10^{-7}. \tag{1.2}$$

Strictly speaking, we would identify $\omega$ in this expression as the distance between single-particle levels of opposite parity, which is smaller than $p^2/2m$ on the average. Usually, however, the matrix element $\langle H_W\rangle$ is also slightly smaller than the rough estimate because of the incomplete overlap of wave functions. The net result is that the approximation $F\sim Gm_\pi^2=2\cdot10^{-7}$ seems reasonable.

Parity breaking in nuclei was first observed in the $^{113}$Cd$(n,\gamma)$ reaction (see the review by Abov and Krupchitskiĭ[14]). The magnitude of the observed asymmetry in the angular distribution of $\gamma$ rays was $-(4.1\pm0.8)\cdot10^{-4}$, or much larger than $F$. This enhancement is due primarily to the high level density in the compound nucleus.[15-17] Following Shapiro,[17] we will refer to this effect as "dynamic enhancement." Let us review its origins. The wave function of any state in a compound nucleus may be expanded in products of single-particle wave functions,

$$\Psi=\sum_{i=1}^{N}a_i\varphi_i, \tag{1.3}$$

where $\varphi_i$ are products of the wave functions of excited particles and holes. The typical number of terms in the sum is determined by the intensity of the residual internucleon interaction. If $\Delta E$ is the scale of this interaction, and $D$ is the distance between levels of the compound nucleus, then the number of terms is $N\sim\Delta E/D$. We recall that $D$ falls off exponentially with increasing number of excited particles. In heavy fissile nuclei ($A\approx240$) we have $D\sim1$ eV. In intermediate nuclei (Cd, Sn, etc.), we have $D\sim10$–$100$ eV. A typical value is $\Delta E\sim\omega\sim1$ MeV ($\omega$ is the distance between single-particle levels), so that we have $N\sim10^4$–$10^6$. Clearly, this strong mixing will make the coefficients $a_i$ comparable in magnitude. By virtue of the normalization condition we can thus write $|a_i|\sim1/\sqrt{N}$.

We turn now to the matrix element of the single-particle operator $H_W$ between two states of a compound nucleus:

$$M=\left\langle\sum_i a_i\varphi_i|H_W|\sum_k b_k\varphi_k\right\rangle=\sum_{i,k}a_i^*b_k\langle i|H_W|k\rangle. \tag{1.4}$$

For each fixed value of $i$, the matrix element $\langle i|H_W|k\rangle$ is nonzero for only a few values of $k$, for which $\varphi_i$ differs from $\varphi_k$ by the state of a single particle. It is natural to assume that the signs of the individual terms in the sum in (1.4) occur at random. We thus have in (1.4) an incoherent sum of $\sim N$ terms each of the order of $\langle H_W\rangle/N$. We can thus write

$$|M|\sim|\langle H_W\rangle|\frac{1}{\sqrt{N}}, \tag{1.5}$$

where $\langle H_W\rangle$ is a typical matrix element between single-particle states. Since the matrix elements of the mixing between different levels of the compound nucleus are of the same order of magnitude, the mixing will be at a maximum between nearest levels. The mixing ratio is given in magnitude by

$$|\alpha|\sim\frac{|M|}{D}\sim\frac{|\langle H_W\rangle|}{\omega}\sqrt{N}=F\sqrt{N}. \tag{1.6}$$

The typical mixing of single-particle levels is $\sim F$; in other words, the enhancement factor is $\sqrt{N}\sim10^3$ in fissile nuclei and $\sim10^2$ in Cd, Sn, etc.

We wish to emphasize that a circumstance of importance to dynamic enhancement is that the intensity of the residual internucleon interaction, which mixes the single-particle levels, is comparable to the distance between single-particle levels of opposite parity. In, for example, the case of a gas of particles which are moving in a common potential but which are not interacting with each other, the spacing of the levels which are mixed will remain equal to the single-particle spacing even at a high level density of the system, and there will be no dynamic enhancement.

TABLE I. The asymmetry parameter in the angular distribution of fission fragments, $a(\times 10^4)$. The asterisks mark values obtained by recalculating the asymmetry coefficient for the emission of fission neutrons.

| $^{233}U$ | $^{235}U$ | $^{239}Pu$ | Reference |
|---|---|---|---|
| — | $1.37 \pm 0.35$ | — | 1 |
| — | $2.2 \pm 1$ | $-4.8 \pm 0.7$ | 2 |
| — | $1.50 \pm 0.44$ | $-4.8 \pm 0.8$ | 3 |
| $2.8 \pm 0.3$ | — | — | 4 |
| — | — | $-7.8 \pm 0.8$* | 5 |
| — | $0.5 \pm 0.3$* | — | 6 |
| $4.83 \pm 0.38$ | | — | 7 |
| — | $0.84 \pm 0.06$ | — | 8 |
| $3.7 \pm 0.6$* | | — | 9 |
| $3.60 \pm 0.34$ | $0.75 \pm 0.12$ | — | 10 |
| $5.28 \pm 0.25$ | | $-6.22 \pm 0.35$ | 11 |
| $4.2 \pm 0.3$ | $1.1 \pm 0.2$ | $-6.7 \pm 0.9$ | 12 |

## 2. PARITY BREAKING IN NUCLEAR FISSION

Vladimirskiĭ and Andreev[18] may be credited with the first suggestion regarding a search for parity breaking in nuclear fission. Specifically, they discussed spontaneous fission. No direct experiment on spontaneous fission has been carried out but experiments on the fission of nuclei by thermal neutrons have proved feasible. In these experiments one measures the asymmetry in the emission directed of the light fragment with respect to the spin of the initial neutron [$W(\theta) = 1 + a \cos\theta$]. The results which have been obtained are listed in Table I.

To explain the mechanism for parity breaking in fission, we must first discuss how the two-particle weak interaction affects the collective (actually macroscopic) motion of the system of heavy fragments. Yet another problem—the problem which was primarily responsible for the cloud of doubt which hung over the first experiment—stems from the large number of final states of the fragments ($N > 10^{10}$). If the sign of the effect depended in a random way on the final state of the system [as it does, for example, in the $(\mathbf{n}, \gamma)$ reaction], then there would be a pronounced suppression of the asymmetry in a real experiment, in which all the final states are detected essentially at once. We must therefore determine why the observed effect does not disappear when an average is taken over the final states of the fragments.

This section of this review is based primarily on the material in Ref. 19. To the best of our knowledge, Danilyan[20] was the first to suggest that parity breaking in a compound nucleus would affect the magnitude of the effect in fission (see also Ref. 21). The relationship between the asymmetry in the fragment emission direction and the mixing of rotational states of the cold nucleus was discussed in Ref. 22.

### a) Fission channels

If the energy is not too high, the fission of a nucleus goes through the following steps: the capture of the neutron and the formation of a hot compound nucleus; the existence of a cold pear-shaped nucleus; and the rupture of the neck connecting the fragments. The entire multitude of final states is formed in this last step. Let us see how the intermediate cold step arises. If we were dealing with the fission of a classical charged drop with an energy just slightly above the fission barrier, then the shape of the drop as it passed the saddle point would essentially be fixed, regardless of the initial conditions, since essentially all the excitation energy is expended on the deformation. The initial conditions in this case determine only the time at which the barrier is reached. In the quantum-mechanical case, in contrast, the energy of the changes in shape is a discrete quantity. During motion near the saddle point, therefore, the nucleus may be in only one or at most a few internal states. Such states, which have definite quantum numbers for all degrees of freedom except the motion across the barrier, are called "fission channels."[23]

Let us put this hypothesis of fission channels in more formal terms. We denote by $|\text{in}\rangle$ the initial state of the system formed by the neutron and the nucleus (with a given asymptotic behavior in the limit $t \to -\infty$), while $\langle\text{out}|$ is a specific final state of the fragments (in this case the asymptotic behavior is specified in the limit $t \to +\infty$). The amplitude for fission is equal to the scalar product $\langle\text{out}|\text{in}\rangle$. This amplitude may be written as a sum over any complete set of intermediate states,

$$\langle\text{out}|\text{in}\rangle = \sum_a \langle\text{out}|a\rangle \langle a|\text{in}\rangle. \tag{2.1}$$

The existence of fission channels means that there is a set $|a\rangle$ for which the sum is determined completely by a few terms. It is clear on physical grounds that this set contains the wave function of the unexcited elongated nucleus and a few of the lowest-lying excited states.

### b) Mixing of states of opposite parity

Let us examine the amplitude $\langle a|\text{in}\rangle$ for a transition from an initial state to a cold intermediate state. Experimentally, the excitation energy of a nucleus is specified very accurately, within $\delta E \sim 0.03$ eV, which is the thermal spread of the neutron energy. In this case it is meaningless to speak in terms of a temporal resolution of the fission process into sequential steps, since the uncertainty relation $\delta t \cdot \delta E \sim \hbar$ tells us that under the condition $\delta E < \Gamma$ the uncertainty in the time satisfies $\delta t > T$, where $T$ is the lifetime of the nucleus. In this situation the fission process can be described satisfactorily by a wave function $\Psi(E)$ corresponding to a fixed energy. Such a wave function incorporates parts corresponding to both the initial compound nucleus and the cold step.

Let us examine the capture of a neutron at an energy $E$ which is approximately equal to the energy of some $s$-wave compound resonance. In the resonance approximation, the wave function of the nucleus after the capture is the same as the wave function of the given compound state [$\Psi(E) \propto \Psi(E_s)$]. It is convenient to single out from the wave function $\Psi$ as a separate term the part ($|a\rangle$) which corresponds to the cold nucleus. This can be done by expanding $\Psi$ in a complete set of states which includes $|a\rangle$ as one of the basis states. We might choose this set to be the products of single-quasiparticle wave functions (where "quasiparticles" are to be understood as excited nucleons, holes, and core vibrations). In these terms the cold step corresponds to the state with maximum deformation, in which only a single degree of freedom—the vibrational degree of freedom—is excited.[1)] We can thus write

$$\Psi_\eta=\sum_{i=1}^{N} a_i\varphi_i^\eta+A_\eta|a\rangle^\eta, \tag{2.2}$$

where $\eta=\pm1$ is the parity. The wave function $\Psi$ includes states $\varphi_i$ in which one, two, etc., quasiparticles are excited. For simplicity we are restricting the discussion at this point to the case of a single fission channel, i.e., to the case in which only a single state $|a\rangle^\eta$ goes into the continuum. The amplitude for a transition to the cold state is $\langle a|\mathrm{in}\rangle\propto A$, and the fission probability is proportional to $|A|^2$. As shown in Sec. 1, the number of terms in $\Psi_\eta$ in a system with strong mixing is $N\sim\Delta E/D\sim10^6$, so that $|A|\sim|a_i|\sim1/\sqrt{N}$. The wave function of opposite parity is of a form similar to that of $\Psi_\eta$:

$$\Psi_{\bar\eta}=\sum_{k=1}^{N} b_k\varphi_k^{\bar\eta}+A_{\bar\eta}|a\rangle^{\bar\eta};\quad |A_{\bar\eta}|\sim|b_k|\sim\frac{1}{\sqrt{N}}. \tag{2.3}$$

The mixing of the nearest $\Psi$ and $\Psi_{\bar\eta}$ levels is dynamically enhanced. The total wave function is

$$\Psi=\Psi_\eta+\alpha\Psi_{\bar\eta}=\sum a_i\varphi_i^\eta+\alpha\sum_k b_k\varphi_k^{\bar\eta}+A_\eta\left(|a\rangle^\eta+\alpha\frac{A_{\bar\eta}}{A_\eta}|a\rangle^{\bar\eta}\right), \tag{2.4}$$

where

$$\alpha=\frac{\langle\Psi_{\bar\eta}|H_W|\Psi_\eta\rangle}{E-E_{\bar\eta}+(i\Gamma_{\bar\eta}/2)}\sim F\sqrt{N}\sim10^{-4}. \tag{2.5}$$

We wish to emphasize that $\Psi_\eta$ and $\Psi_{\bar\eta}$ correspond to quasistationary states in the continuum. It is for this reason that the energy denominator in (2.5) is of the form $E-E_{\bar\eta}+(i\Gamma_{\bar\eta}/2)$ ($E$ is the energy at which the capture occurs) rather than $E_\eta-E_{\bar\eta}$.

Let us examine the cold step of the fission. In this step the nucleus is a pear-shaped top. The spectrum of such a system may be described as follows in the adiabatic approximation.[23] For a fixed internal state $|aK\rangle$ ($K$ is the projection of the total angular momentum $J$ onto the axis of the top), there is a band of rotational states. If $K\neq0$, then there are two rotational levels of opposite parity ($\eta=\pm1$) for each value of $J$—an effect analogous to the $\Lambda$-doubling effect in heteronuclear molecules.[2)] (In the case of a quadrupole deformation, as in homonuclear molecules, one of the doublet levels would be forbidden by the requirements of the statistics.) For $K=0$ the parity is related to $J$ unambiguously: $\eta=(-1)^J\times$ (the internal parity), so that for each $J$ there is only a single level. In the adiabatic approximation the wave function of the rotating nucleus may be written

$$\left.\begin{aligned}
&|aK\rangle_M^{\eta}=\frac{1}{\sqrt2}\left[|aKJM\rangle+\eta|a\bar KJM\rangle\right],\\
&|aKJM\rangle=\sqrt{\frac{2J+1}{4\pi}}\,D^J_{MK}(\varphi,\theta,0)|aK\rangle,\\
&|a\bar KJM\rangle=(-1)^{J+K}|a,-KJM\rangle.
\end{aligned}\right\} \tag{2.6}$$

[1)]The appearance of the cold step as a separate term in the total wave function of the nucleus is especially graphic in an analysis of nuclear fission by perturbation theory in which the interaction of the excited nucleons with core vibrations is treated as a perturbation (see the Appendix).

[2)]This is precisely the rotational structure which has been observed in the fission of $^{230}$Th by neutrons.[24]

We have chosen the phase of the wave function such that at $K=0$ the angular part is equal to $Y_{JM}(\theta,\varphi)$. This definition differs from that adopted in Ref. 23 by the absence of a factor $i^{(1-\eta)/2}$.

For simplicity we consider the case in which the neutron is captured into a resonance with a fixed value of $J$, and the fission goes through a single channel with $K\neq0$. The part of the wave function in (2.4) corresponding to the cold step is

$$\begin{aligned}
|\widetilde{aK}\rangle_M^{\eta}&=|aK\rangle_M^\eta+\beta|aK\rangle_M^{\bar\eta}\\
&=\frac{1}{\sqrt2}\left[(1+\beta)|aKJM\rangle+\eta(1-\beta)|a\bar KJM\rangle\right],
\end{aligned} \tag{2.7a}$$

$$\beta=\sum_\nu\alpha_\nu\frac{A_{\bar\eta}(\nu)}{A_\eta}. \tag{2.7b}$$

Equations (2.7) differ from (2.4) in that several admixed levels, rather than a single one, are taken into account.

### c) Angular distribution of fragments

Comparing (2.7) with (2.1), we see that we have actually calculated the wave function $\sum_a|a\rangle\langle a|\mathrm{in}\rangle$. Now we need to project this wave function onto the final state,

$$|\mathrm{out}\rangle=|fK\rangle_M^\eta=\frac{1}{\sqrt2}\left(|fKJM\rangle+\eta|f\bar KJM\rangle\right). \tag{2.8}$$

The main and admixed wave functions in (2.7a) correspond to the same internal state of the nucleus and differ only in the macroscopic rotational motion; specifically, angular momenta of different parity appear in the expansion of this state in terms of the orbital angular momenta in the states $\eta$ and $\bar\eta$. The amplitudes for fission from the states $|aK\rangle^\eta$ and $|aK\rangle^{\bar\eta}$ to any specific internal state of the fragments are therefore equal, and the wave function in terms of angular variables is also of the form in (2.7) in the limit $r\to\infty$. Specifically,

$$\begin{aligned}
\sum_a|f\rangle\langle f|a\rangle\langle a|\mathrm{in}\rangle&\propto|fKJM\rangle\langle fKJM|aKJM\rangle(1+\beta)\\
&+\eta|f\bar KJM\rangle\langle f\bar KJM|a\bar KJM\rangle(1-\beta)\propto|fK\rangle_M^\eta+\beta|fK\rangle_M^{\bar\eta}.
\end{aligned} \tag{2.9}$$

Here we have used conservation of the quantum number $K$ in the course of rupture and the equality of amplitudes $\langle fKJM|aKJM\rangle=\langle f\bar KJM|a\bar KJM\rangle$ which follows from parity conservation in rupture. Squaring (2.9), and substituting $|fK\rangle_{JM}^\eta$, we find from (2.8) and (2.6) the angular distribution of the fragments:

$$W_{JM}(\theta)\propto|D^J_{MK}|^2(1+\gamma)+|D^J_{M,-K}|^2(1-\gamma), \tag{2.10}$$

$$\gamma=2\,\mathrm{Re}\,\beta=2\,\mathrm{Re}\left(\sum_\nu\sqrt{\frac{\Gamma_\nu}{\Gamma}}\frac{\langle\bar\eta\nu|H_W|\eta\rangle}{E-E_\nu+(i\Gamma_\nu/2)}e^{i(\varphi-\varphi_\nu)}\right). \tag{2.11}$$

Here we have used $A_\eta\propto\sqrt{\Gamma}e^{i\varphi}$ and $A_{\bar\eta}(\nu)\propto\sqrt{\Gamma_\nu}e^{i\varphi_\nu}$, where $\Gamma$, $\Gamma_\nu$, $\varphi$, and $\varphi_\nu$ are the fission widths and the phase shifts of the transition to the cold state from the corresponding levels of the compound nucleus. The question of phase shifts will be discussed in more detail in the Appendix. For the above-barrier fission with which we are concerned here, we evidently have $\Gamma_\nu\sim\Gamma$; this result has been confirmed experimentally. We thus have

$$\gamma \sim \beta \sim \frac{F}{\sqrt{N}} \sim 10^{-4}. \tag{2.12}$$

For fission by tunneling, there may be situations in which $\Gamma_\nu$ is larger than $\Gamma$; the corresponding enhancement of the effect has been discussed elsewhere.[18,25]

In the experiments of Refs. 1–12, polarized neutrons caused the fission of unpolarized nuclei. The angular distribution for this case is

$$\begin{gathered} W(\theta) \propto \sum_M |C^{JM}_{IM-\frac{1}{2},\frac{1}{2}\frac{1}{2}}|^2 W_{JM}(\theta) \sim 1 + a\cos\theta, \\ a = \gamma \frac{K}{I+(1/2)} (-1)^{J-I-(1/2)}, \end{gathered} \tag{2.13}$$

where $I$ is the spin of the target nucleus, and $C^{JM}_{IM-1/2,1/2\,1/2}$ is the Clebsch-Gordan coefficient.

We wish to emphasize that the effect is a consequence of the orientation of the nucleus before rupture. In the state (2.7) there is an average orientation of the nucleus along $\mathbf{J}$. The $\mathbf{Jn}$ correlation ($\mathbf{n}$ is the direction of the axis of the nucleus) is not only $P$-odd but also $T$-odd. This correlation can therefore arise only as a result of a finite lifetime of the nucleus. In this case the spectrum of the system is continuous, rather than discrete, and thus degenerate: A state with a given energy may correspond to both an outgoing wave and an incoming wave. The outgoing wave, in which we are interested, is not an eigenfunction of the time-reversal operator ($T$ reversal transforms this wave into an incoming wave), so that $T$-odd correlations are not forbidden.

There is another way to explain the origin of the $\mathbf{Jn}$ correlation. In an unstable nucleus, the average value of the momentum of the fragment along the direction of the axis, $\mathbf{n}$ is not zero ($\mathbf{p} \propto \Gamma\mathbf{n}$). The $\mathbf{Jp}$ correlation, on the other hand, is $T$-even.

Since the effect stems from the finite lifetime of the nucleus, it must vanish in the limit $\Gamma \to 0$. That this is in fact the case can be seen easily and directly. For our choice of wave-function phases, the weak-interaction matrix element is purely imaginary, so that there is no $\mathbf{Jn}$ correlation in the state (2.7) with $\Gamma_\nu = 0$ and $\cos(\varphi - \varphi_\nu) = \pm 1$, and the effect vanishes[3)] ($\gamma = a = 0$). A question which naturally arises is whether a phase shift between states of opposite parity might appear in the free motion of the fragments away from the rupture point to infinity. There is such a phase shift, of course, but it is a consequence of the violation of the adiabatic approximation, and it is negligibly small. Indeed, the reason for the appearance of a phase shift during the free motion is the difference between the centrifugal potential energies for the radial motion. The adiabatic approximation with respect to rotation, however, means precisely that the centrifugal energy is ignored. If the justification for the use of the adiabatic approximation for a deformed nucleus is a rotational energy which is small in comparison with the energy of the internal excitation, then in the free motion the rotational energy ($E_{\text{rot}} \sim \hbar^2/MR^2 \le 10$ keV) is small in comparison with the kinetic energy of the fragments ($E_{\text{kin}} > 10$ MeV). The phase shift of the free motion is thus $\Delta\varphi \sim \sqrt{E_{\text{rot}}/E_{\text{kin}}} \sim 10^{-2}$. To avoid confusion we emphasize that we are here talking about motion of particles of finite size. In the motion of a point particle away from the origin of coordinates, the phase shift of the free motion in states of opposite parity (with orbital angular momenta $l$ and $l+1$) is always $\pi/2$, because of a singularity of the centrifugal energy $l(l+1)/2mr^2$ in the limit $r \to 0$.

As we mentioned earlier, if a parity-breaking effect is to occur, it is important that the energy spread of the neutrons be smaller than the distance between the levels of the compound nucleus and also smaller than the width of these levels ($\delta E < D$, $\Gamma$). It is in this case that it is sufficient to deal with the wave function of a quasistationary state with a given energy, rather than the time evolution of the fission. In the opposite case $\delta E \gg D \gtrsim \Gamma$ a temporal description is possible. In that case, however, the initial hot step is "forgotten" upon the transition to the cold state, and the effect caused by the weak interaction in the hot step is suppressed by a factor $\sqrt{\delta E/D}$.

### d) Overlap of neutron resonances. Dependence of the effect on the final state of the fragments

In a real situation there may be an interference among several incoming neutron resonances. If all have the same $J$, then Eq. (2.10) is again applicable, while expression (2.11) for $\gamma$ will be modified in an obvious manner. A more interesting effect arises if resonances with different $J$ overlap. We denote by $T(+)$ the amplitude for capture to a resonance with $J_+ = I + (\frac{1}{2})$ and by $T(-)$ the amplitude for capture to a resonance with $J_- = I - (\frac{1}{2})$. The wave function in the cold state is then

$$\begin{gathered} C^{J_+M}_{IM-\frac{1}{2},\frac{1}{2}\frac{1}{2}} u_+ |\overline{aK}\rangle^{\eta}_{J_+M} + C^{J_-M}_{IM,-\frac{1}{2},\frac{1}{2}\frac{1}{2}} u_- |\overline{aK}\rangle^{\eta}_{J_-M}, \\ u_\pm = \frac{T(\pm) A_\eta(\pm)}{E - E_\pm + (i\Gamma_\pm/2)}. \end{gathered} \tag{2.14}$$

The asymmetry coefficient for the fission of unpolarized nuclei by polarized neutrons is

$$\begin{gathered} a = \frac{2}{|v_+|^2 + |v_-|^2} \left\{ \frac{K}{I+(1/2)} (|v_+|^2 \operatorname{Re}\beta_+ - |v_-|^2 \operatorname{Re}\beta_-) \right. \\ \left. - \frac{\sqrt{[I+(1/2)]^2 - K^2}}{I+(1/2)} \operatorname{Re}[v_+ v_-^* (\beta_+ + \beta_-^*)] \right\}, \\ v_\pm = \sqrt{2J_\pm + 1}\, u_\pm. \end{gathered} \tag{2.15}$$

If $u_+$ or $u_-$ is zero, the result is naturally the same as (2.13). We note that, in contrast with the first two terms, the interference term does not fall off with increasing spin ($I$) of the initial nucleus.

If $K = 0$, the cold nucleus has no rotational levels of different parity and identical $J$. Because of the interference of neutron resonances with different $J$, however, the effect does not vanish even at $K = 0$. To demonstrate this assertion, we consider the par-

[3)]In the problem under consideration here, this circumstance does not lead to any significant decrease in the effect, if only because we have $E - E_\nu \sim \Gamma_\nu$. Furthermore, there is no reason to assume that the difference between the nonresonant phase shifts, $\varphi - \varphi_\nu$, is small.

ticular case of the reaction $^{239}Pu(n,f)$ $(I=\frac{1}{2})$. The thermal neutrons are captured into the $|0^+\rangle$ and $|1^+\rangle$ resonances. With $K=0$, the parity of the cold nucleus is $\eta=(-1)^J$, so that fission is allowed from the $|0^+\rangle$ resonance, while fission from the $|1^+\rangle$ resonance through a channel with $K=0$ is forbidden. The weak interaction, however, mixes the $|1^-\rangle$ compound state (from which fission is possible) with the $|1^+\rangle$ resonance. The wave function $Y_{1M}(\theta,\varphi)$ is therefore mixed with the angular wave function of the nucleus in the cold step, $Y_{00}(\theta,\varphi)$. It is not difficult to see that in this case the angular asymmetry is

$$a(K=0)=-2\sqrt{3}\,\mathrm{Re}\left\{\frac{T(1^+)}{T(0^+)}\frac{A(1^-)}{A(0^+)}\frac{\langle 1^+|H_W|1^-\rangle}{E-E(1^-)+[i\Gamma(1^-)/2]}\times\frac{E-E(0^+)+[i\Gamma(0^+)/2]}{E-E(1^+)+[i\Gamma(1^+)/2]}\right\}. \quad (2.16)$$

Let us examine a possible dependence of the effect on the final states of the fragments. We should distinguish between two possibilities: rapid fluctuations of the effect (the effect changes noticeably from one specific final state of the system, $|f\rangle$, to another[4)]) and a slow change, i.e., a smooth dependence on total (macroscopic) characteristics, e.g., on $\xi=(A_1-A_2)/(A_1+A_2)$, where $A_1$ and $A_2$ are the fragment masses. Up to this point we have been discussing the case of a single fission channel. As we saw in the preceding subsection, there is no dependence of the asymmetry on the final state. This assertion, of course, holds only within the range of applicability of the adiabatic approximation for the wave function of the cold nucleus [Eq. (2.6)]. The adiabatic approximation is obviously violated in the limits $\xi\to 1$ and $\xi\to 0$. In the limit $\xi\to 1$ the mass of one of the fragments is small, and the characteristic rotational energy becomes comparable to the energy of the internal excitations. In the limit $\xi\to 0$, on the other hand, a small excess of nucleons may tunnel from one fragment to another, and this event would also cause pronounced separations of rotational levels of opposite parity.

We turn now to the more realistic case in which the fission goes through several channels. If these channels have different values of $K$, they do not interfere with each other (we are assuming that $K$ is conserved during rupture), and we can write $a(f)=\sum_i W_i(f)a_i$, where $W_i(f)$ is the relative probability for fission through the $|i\rangle$ channel to the state $|f\rangle(\sum_i W_i(f)=1)$. We see that rapid fluctuations have appeared in the asymmetry because of $W_i(f)$. In actual experiment, an average is always taken over some interval of final states; in other words, one actually measures

$$a=\overline{a(f)}=\sum_i\overline{W_i(f)}\,a_i.$$

Finally, if there are several channels with identical $K$, they may interfere, causing a complicated dependence on $|f\rangle$. It seems natural, however, that the interference would disappear after an average is taken over a certain interval of final states, and we would again have $a(\xi)=\sum_i\overline{W_i(f)}\,a_i=\sum_i W_i(\xi)a_i$ for the observable quantity. We see that the effect may depend on $\xi$ even in the adiabatic region if the fragment mass distributions are different for different channels. We must add that it would be very difficult to imagine this difference to be large, so that the experimental absence of a dependence of the asymmetry on the fragment masses[11] seems quite plausible.

We wish to emphasize again that the angular asymmetry arises in the cold step, before the nucleus ruptures to form the fragments (see Subsec. 2c). It is for this reason that the effect is not suppressed when an average is taken over the final states of the fragments.

## 3. P-EVEN ANGULAR CORRELATIONS IN THE FISSION OF UNPOLARIZED NUCLEI BY SLOW NEUTRONS

From both the experimental and theoretical standpoints, the $P$-even angular correlations are intimately related to parity breaking in fission.[10,11,26] Furthermore, it is fair to say that the observation of these correlations have served as a strong argument in favor of the parity-breaking mechanism discussed above. We therefore believe it is necessary to take up the question of $P$-even correlations here. The correlations involved are the $\mathbf{pk}$ and $\mathbf{p}[\mathbf{k}\sigma]$ correlations ($\sigma$ and $\mathbf{k}$ are the spin and momentum of the neutron, and $\mathbf{p}$ is the momentum of the light fragment). From the standpoint of the formation mechanism, these correlations differ from the $P$-odd effect only in that in the even case the mixing of the compound-nucleus levels results not from the weak interaction but from an overlap of $s$-wave and $p$-wave neutron resonances.

We expand the wave function of the incident neutron in terms of states with a definite angular momentum, $|ljj_z\rangle$:

$$e^{i\mathbf{kr}}X_\alpha=4\pi\sum_{l,m}i^l j_l(kr)Y^*_{lm}(\mathbf{n}_k)Y_{lm}(\mathbf{n})X_\alpha=4\pi\Big(Y^*_{00}(\mathbf{n}_k)\Big|0,\frac{1}{2},\alpha\Big\rangle+i\frac{kr}{3}\sum_m Y^*_{1m}(\mathbf{n}_k)C^{jj_z}_{1m\frac{1}{2}\alpha}|1,j,j_z\rangle+\ldots\Big); \quad (3.1)$$

where $X_\alpha$ is the spin wave function with spin projection $\alpha$, $j_l(kr)$ is the spherical Bessel function, $\mathbf{n}_k=\mathbf{k}/k$, and $\mathbf{n}=\mathbf{r}/r$. For slow neutrons, the terms with $l>1$ are inconsequential. Neutrons from $s$ and $p$ waves are captured into compound-nucleus states of opposite parity. The wave function of the nucleus after capture of a neutron from the state (3.1) may be written as

$$\sum_s\frac{T_s}{E-E_s+(i\Gamma_s/2)}Y^*_{00}(\mathbf{n}_k)C^{J_sJ_{sz}}_{II_z\frac{1}{2}\alpha}|\eta,J_s,J_{sz}\rangle+\sum_p\sum_{j,m}\frac{iT_p(j)}{E-E_p+(i\Gamma_p/2)}Y^*_{1m}(\mathbf{n}_k)C^{J_pJ_{pz}}_{II_zjj_z}C^{jj_z}_{1m\frac{1}{2}\alpha}|\bar{\eta},J_p,J_{pz}\rangle; \quad (3.2)$$

where $\eta$ is the parity of the target nucleus, and $\bar{\eta}\equiv-\eta$. The Clebsch-Gordan coefficient $C^{JJ_z}_{II_zjj_z}$ arises when the spin of the target nucleus, $I$, is combined

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4)We wish to thank V. M. Strutinskiĭ and Yu. V. Petrov for calling our attention to the question of fluctuations of the effect.

with the angular momentum of the neutron. The amplitudes $T_s$ and $T_p(j)$ correspond to the capture of a neutron from the states $|l=0, j=\frac{1}{2}, \alpha\rangle$ and $|l=1, j, j_z\rangle$ into the compound-nucleus states $|\eta, J, J_z\rangle$ with a given total angular momentum $J$. Obviously, we have $T_p/T_s \sim kR$ where $R$ is the nuclear radius. Using expressions (2.2) and (2.3) for the states $|\eta, J_s, J_{sz}\rangle$ and $|\bar{\eta}, J_p, J_{pz}\rangle$, we can single out from the wave function in (3.2) a part which corresponds to the cold nucleus:

$$\sum_s u_s Y^*_{00}(\mathbf{n}_k) C^{J_s J_{sz}}_{II_z \frac{1}{2}\alpha} |aK\rangle^{\eta}_{J_s J_{sz}} + \sum_p \sum_{j,m} u_{pj} Y^*_{1m}(\mathbf{n}_k) C^{J_p J_{pz}}_{II_z jj_z} C^{jj_z}_{1m\frac{1}{2}\alpha} |a, K\rangle^{\bar{\eta}}_{J_p J_{pz}}, \qquad (3.3)$$

where

$$u_s = \frac{T_s A_s}{E - E_s + (i\Gamma_s/2)}, \quad u_{pj} = \frac{iT_p(j) A_p}{E - E_p + (i\Gamma_p/2)};$$

and $A_s$ and $A_p$ are the corresponding amplitudes for a transition to the cold step. As we mentioned earlier, the angular part of the wave function will also have the form of (3.3), with the wave functions $|a, K\rangle^{\eta}_{JJ_z}$ from (2.6), in the limit $r \to \infty$. To find the angular distribution it is thus sufficient to average the square modulus of wave function (3.3) over the spin projections of the target nucleus, $I_z$. After some straightforward but lengthy calculations, we find the following expression for the angular distribution:

$$W(\mathbf{n}_p) \propto \sum_{s,s'} (2J_s + 1) u_s u^*_{s'} \delta_{J_s J_{s'}} + \sum_{s,p,j} Q(J_s, J_p, j, K, I) \operatorname{Re}\{u_s u^*_{pj} (\mathbf{n}_k \mathbf{n}_p - i\beta_j \mathbf{n}_p [\mathbf{n}_k \sigma])\}, \qquad (3.4)$$

where $\mathbf{n}_p = \mathbf{p}/p$ is the emission direction of the light fragment, $|\sigma| = 1$, and

$$\beta_j = \begin{cases} 1, & j = \frac{1}{2}, \\ -\frac{1}{2}, & j = \frac{3}{2}, \end{cases}$$

$$Q(J_s, J_p, j, K, I) = 2\sqrt{3}(2J_s + 1)(2J_p + 1)\sqrt{2j+1} \times (-1)^{1+j-I-K} \begin{Bmatrix} \frac{1}{2} & 1 & j \\ J_p & I & J_s \end{Bmatrix} \begin{pmatrix} J_s & J_p & 1 \\ K & -K & 0 \end{pmatrix}. \qquad (3.5)$$

Let us consider the simplest case, in which there is only a single $s$-wave resonance and a single $p$-wave resonance near the thermal region. It is convenient to express the ratio of capture amplitudes, $T_p/T_s$, and the ratio of fission amplitudes, $A_p/A_s$, in terms of the ratios of the corresponding widths:

$$\frac{T_p}{T_s} = \sqrt{\frac{\Gamma^{(p)}_n}{\Gamma^{(s)}_n}} \exp[i(\varphi^{(p)}_n - \varphi^{(s)}_n)] = \pm\sqrt{\frac{\Gamma^{(p)}_n}{\Gamma^{(s)}_n}},$$
$$\frac{A_p}{A_s} = \sqrt{\frac{\Gamma^{(p)}_f}{\Gamma^{(s)}_f}} \exp[i(\varphi^{(p)}_f - \varphi^{(s)}_f)]. \qquad (3.6)$$

Here we have used $\varphi^{(p)}_n - \varphi^{(s)}_n = m\pi$, where $m$ is an integer, as the difference between the capture phase shifts for slow neutrons ($kR \ll 1$). [The phase shift of the free motion of a $p$ wave, $e^{i\pi/2} = i$, is not incorporated in the definition of $T_p$ and is singled out explicitly in (3.2) and (3.3).] The angular distribution is

TABLE II. The factor $Q(J_s, J_p, j, K, I)/(2J_s + 1)$.

| | $J_s$ | $J_p$ | $K=0$ | | $K=1$ | | $K=2$ | |
|---|---|---|---|---|---|---|---|---|
| | | | $j=1/2$ | $j=3/2$ | $j=1/2$ | $j=3/2$ | $j=1/2$ | $j=3/2$ |
| $^{239}$Pu, $I=1/2$ | 0 | 1 | 3.46 | 4.90 | — | — | — | — |
| | 1 | 0 | 1.15 | — | — | — | — | — |
| | 1 | 1 | 0 | 0 | −2.00 | 1.41 | — | — |
| | 1 | 2 | — | 3.65 | — | 3.16 | — | — |
| $^{233}$U, $I=5/2$ | 2 | 1 | — | −1.70 | — | −1.47 | — | — |
| | 2 | 2 | 0 | 0 | 0.67 | 1.25 | 1.33 | 2.49 |
| | 2 | 3 | 2.37 | 2.12 | 2.23 | 2.00 | 1.76 | 1.58 |
| | 3 | 2 | 1.69 | −0.90 | 1.59 | −0.85 | 1.26 | −0.67 |
| | 3 | 3 | 0 | 0 | −0.67 | 0.75 | −1.33 | 1.49 |
| | 3 | 4 | — | 2.97 | — | 2.87 | — | 2.57 |
| $^{235}$U, $I=7/2$ | 3 | 2 | — | −1.92 | — | −1.81 | — | −1.43 |
| | 3 | 3 | 0 | 0 | 0.50 | 0.87 | 1.00 | 1.73 |
| | 3 | 4 | 2.27 | 1.92 | 2.20 | 1.86 | 1.96 | 1.66 |
| | 4 | 3 | 1.76 | −1.02 | 1.71 | −0.99 | 1.53 | −0.88 |
| | 4 | 4 | 0 | 0 | −0.50 | 0.59 | −1.00 | 1.18 |
| | 4 | 5 | — | 2.85 | — | 2.80 | — | 2.62 |

$$W(\mathbf{n}_p) = 1 + \sum_{j=\frac{1}{2},\frac{3}{2}} b(\mathbf{n}_k \mathbf{n}_p \cos\varphi - \mathbf{n}_p [\mathbf{n}_k \sigma] \beta_j \sin\varphi),$$
$$b = \frac{Q(J_s, J_p, j, K, I)}{2J_s + 1} \sqrt{\frac{\Gamma^{(pj)}_n \Gamma^{(p)}_f}{\Gamma^{(s)}_n \Gamma^{(s)}_f}} \left| \frac{E - E_s + (i\Gamma_s/2)}{E - E_p + (i\Gamma_p/2)} \right|, \qquad (3.7)$$
$$\varphi = \frac{\pi}{2} + m\pi - \varphi^{(p)}_f - \varphi^{(s)}_f + \arg \frac{E - E_s + (i\Gamma_s/2)}{E - E_p + (i\Gamma_p/2)}.$$

The reason why the $\mathbf{n}_k \mathbf{n}_p$ correlation is proportional to $\cos\varphi$, while the $\mathbf{n}_p[\mathbf{n}_k\sigma]$ correlation is proportional to $\sin\varphi$, lies in the different $T$ parities of these correlations (the $\mathbf{n}_p[\mathbf{n}_k\sigma]$ correlation is $T$-odd). Table II lists numerical values of the coefficient $Q/(2J_s + 1)$ for angular momenta of $^{239}$Pu, $^{233}$U, and $^{235}$U. We see that $Q/(2J_s + 1) \sim 1$. The phase shift $\varphi$ in the fissile nuclei is not small ($|\cos\varphi| \sim |\sin\varphi| \sim 1$), if only because $E - E_s \sim \Gamma$. The factor $|E - E_s + (i\Gamma_s/2)||E - E_p + (i\Gamma_p/2)|^{-1}$, like the phase shifts, does not have any important effect on the parity-breaking effect for thermal neutrons, since $|E - E_s| \sim |E - E_p|$. A possible exception to this assertion is $^{239}$Pu, for which the distance from the thermal region to the $s$-wave resonance is nearly an order of magnitude smaller than the average distance between resonances. Since the probabilities for above-barrier fission from compound levels of different parities are of the same order of magnitude ($\Gamma^{(p)}_f \sim \Gamma^{(s)}_f$), we find the following estimate of the asymmetry for thermal neutrons:

$$b \sim \sqrt{\frac{\Gamma^{(p)}_n}{\Gamma^{(s)}_n}} \sim kR \approx 3 \cdot 10^{-4}. \qquad (3.8)$$

This value is consistent with the experimental data available:

$$10^4 \sum_i b\beta_j \sin\varphi = \begin{cases} \left.\begin{matrix} -3.24 \pm 0.33^{10} \\ -6.43 \pm 0.51^{11} \end{matrix}\right\}, & ^{233}\text{U}, \\ 1.65 \pm 0.11^{10}, & ^{235}\text{U}, \\ 1.25 \pm 0.29^{11}, & ^{239}\text{Pu}. \end{cases} \qquad (3.9)$$

Unfortunately, it is not possible to make a more accurate comparison with experiment because the parameters of the $p$-wave resonances and the phase shifts are not known. Even order-of-magnitude estimates, however, show that the model is valid. In the first place, the hot step is not "forgotten" in the system; in other words, the mixing of levels of different parity is in fact faithfully transferred from the hot step to the cold step. Second, there is no averaging of the effect because of the very large number of final states of the fragments; in other words, the mass and angular asymmetries are shaped before the rupture of the nucleus and the formation of the fragments. Since the $P$-odd effect in fission arises in a corresponding way, according to our arguments, we may assert that the observation of $P$-even or $P$-odd correlations at the predicted level would be evidence for the validity of our model for the shaping of the angular distributions during fission.

## 4. PARITY BREAKING IN NEUTRON OPTICS

In this section we will discuss some new possibilities for studying parity breaking in interactions of neutrons with nuclei; the effects involved here are analogous to effects which have been observed in ordinary optics (Ref. 27; see also the review in Ref. 28). The first suggestions regarding the observation of parity breaking in neutron optics can be credited to Michel[29] and (later) Stodolsky.[30] The idea was to measure the rotation of the spin of a neutron around its momentum direction in a medium. Another effect was also discussed: the appearance of a longitudinal polarization of originally unpolarized neutrons. The rotation angle $\psi$ and the degree of longitudinal polarization, $a$, over a neutron mean free path were predicted to be $\psi \sim 10^{-6}-10^{-8}$ rad and $a \sim 10^{-8}\sqrt{E}$, where $E$ is in electron volts. In 1976, Forte[31] noted that the effect was enhanced near a single-particle $p$-wave resonance. A specific suggestion was to measure the neutron-spin rotation angle in $^{124}$Sn, where there is a $p$-wave resonance at 62 eV with a relatively large single-particle component. According to Forte,[32] there would be a rotation angle $\psi \sim 10^{-3}-10^{-4}$ rad at the optimum neutron energy, and there would be a longitudinal polarization $a \sim 10^{-5}-10^{-6}$ (over a mean free path). For thermal neutrons, the prediction was $\psi_t \lesssim 5\cdot 10^{-6}$ rad/cm (Ref. 32).

In Refs. 29–32, as in subsequent papers by other authors,[33–39] the spin rotation in question was a rotation caused by the scattering of the neutron by the $P$-odd potential of the nucleus; in other words, the nucleus was treated as a particle having no internal degrees of freedom. As we will see below, in heavy and intermediate nuclei there is also a substantial rotation caused by another mechanism, involving virtual excitation of the nucleus to a compound state (Ref. 40; see also Ref. 41).

### a) Resonance effects

Let us examine the capture of a neutron of energy $E$ into a $p$-wave resonance. After the capture, the nucleus goes to some compound state with angular momentum $J$ and parity $\eta$. Actually, because of the weak interaction between nucleons this state is a superposition of states of different parities:

$$\begin{gathered}\psi(E)=|J^{\eta}\rangle+i\sum_{\nu}\varepsilon_{\nu}(E)\,|J^{\bar{\eta}},\ \nu\rangle,\\ i\varepsilon_{\nu}(E)\equiv\alpha_{\nu}=\frac{\langle J^{\bar{\eta}},\ \nu\,|H_W|\,J^{\eta}\rangle}{E-E_{\nu}+(i\Gamma_{\nu}/2)}.\end{gathered}\tag{4.1}$$

We have singled out an "$i$" ($=\sqrt{-1}$) in the expression for the mixing ratio, since the ordinary definition of the angular wave functions would make the matrix element $H_W$ purely imaginary. Because of the dynamic enhancement, the mixing of nearest levels is $\varepsilon \sim F\sqrt{N} \sim 10^{-4}-10^{-5}$ [see (1.6)].

The capture into the state (4.1) comes from both the $p$ and $s$ states of the neutron. Let us expand the wave function of a slow neutron with momentum $k$ and helicity $\pm\frac{1}{2}$ in terms of states with a definite angular momentum $|ljj_z\rangle$ (the $z$ axis is along $k$):

$$\begin{aligned}e^{i\mathbf{kr}}\chi_{\pm}&\approx(1+i\mathbf{kr})\,\chi_{\pm}=\sqrt{4\pi}\left[Y_{00}(\mathbf{n})+i\frac{kr}{\sqrt{3}}Y_{10}(\mathbf{n})\right]\chi_{\pm}\\ &=\sqrt{4\pi}\left(\left|0,\frac{1}{2},\pm\frac{1}{2}\right\rangle\mp i\frac{kr}{3}\left|1,\frac{1}{2},\pm\frac{1}{2}\right\rangle\right.\\ &\qquad\left.+i\frac{kr\sqrt{2}}{3}\left|1,\frac{3}{2},\pm\frac{1}{2}\right\rangle\right),\end{aligned}\tag{4.2}$$

where $\chi_{\pm}$ is the spin wave function. The amplitude for neutron capture from state (4.2) into compound-nucleus state (4.1) is

$$\begin{aligned}T=&\left[\pm T_p\left(j=\frac{1}{2}\right)+\sum_{\nu}\varepsilon_{\nu}T_{s\nu}\right]C^{JJ_z}_{II_z\frac{1}{2},\pm\frac{1}{2}}\\ &+T_p\left(j=\frac{3}{2}\right)C^{JJ_z}_{II_z\frac{3}{2},\pm\frac{1}{2}},\end{aligned}\tag{4.3}$$

where $I$ and $J$ are the angular momenta of the nucleus before and after the capture, and $T_s$ and $T_p(j)$ are the scalar amplitudes for capture from the $s$ and $p$ waves. The "$i$" in (4.1) cancels out with the difference between the free-motion phase shifts of the $s$ and $p$ waves. Working from (4.3), we easily find that the capture amplitudes for different values of $j$ do not interfere if the target is unpolarized. The interference of amplitudes of different parity with $j=\frac{1}{2}$, on the other hand, leads to a difference between the cross sections for the absorption of neutrons with helicities $\pm\frac{1}{2}$,

$$\begin{gathered}\sigma^a_{\pm}=\sigma^a[1\pm P(k)],\quad P(k)=\frac{P_0(k)\,\Gamma_n^{(p\frac{1}{2})}}{\Gamma_n^{(p\frac{1}{2})}+\Gamma_n^{(p\frac{3}{2})}},\\ P_0(k)=2\,\mathrm{Re}\left(\sum_{\nu}\varepsilon_{\nu}(E)\frac{T_{s\nu}}{T_p(j=1/2)}\right)\\ =2\sum_{\nu}\varepsilon_{\nu}(E)\sqrt{\frac{\Gamma_n^{(s\nu)}(k)}{\Gamma_n^{(p\frac{1}{2})}(k)}}\cos\left(\varphi^{(s\nu)}-\varphi^{(p\frac{1}{2})}\right);\end{gathered}\tag{4.4}$$

where $\Gamma_n^{(p)}(k)=\Gamma_n^{(p1/2)}(k)+\Gamma_n^{(p3/2)}(k)$ and $\Gamma_n^{(s\nu)}(k)$ are the neutron widths of the states $|J^{\eta}\rangle$ and $|J,\bar{\eta}\,\nu\rangle$, corrected to the energy of the incident neutron:

$$\Gamma_n^{(pj)}(k)=\Gamma_n^{(pj)}\left(\frac{k}{k_p}\right)^3,\quad\Gamma_n^{(s)}(k)=\Gamma_n^{(s)}\frac{k}{k_s}.\tag{4.5}$$

Here $k_p$ and $k_s$ are the momenta corresponding to the resonances, and $\varphi^{(p)}$ and $\varphi^{(s)}$ are the capture phase shifts. For slow neutrons ($kR\ll 1$, where $R$ is the nuclear radius), we of course have $\cos(\varphi^{(s)}-\varphi^{(p)})=\pm 1$. The case in which we are interested is that in which the distance to the $p$-wave resonance is much smaller than that to the nearest $s$-wave resonances. In this case we can easily single out the explicit functional form of $P(k)$ and $P_0(k)$:

$$P(k)=P\frac{k_p}{k},\quad P_0(k)=P_0\frac{k_p}{k}.\tag{4.6}$$

For low-lying resonances (with $E_p=1-10$ eV) we have $\sqrt{\Gamma_n^{(s)}(k_p)/\Gamma_n^{(p)}(k_p)}\sim 1/k_pR\sim 10^2-10^3$ and $P\lesssim P_0\sim 10^{-1}-10^{-2}$.

Noting that $\sigma^a$ in (4.4) has the standard Breit-Wigner form, we can easily convert from the absorption cross section to a refractive index:

$$n_{\pm}=n_0-\frac{\pi N g\Gamma_{\mathrm{n}}^{(p)}}{k_p^3}\left(1\pm P\frac{k_p}{k}\right)\frac{1}{E-E_p+(i\Gamma/2)},\qquad g=\frac{2J+1}{2(2I+1)};\tag{4.7}$$

where $N$ is the density of target atoms, $n_0$ is the nonresonant part of the refractive index, and $\Gamma$ is the total width of the $p$-wave resonance. We have ignored the Doppler broadening of the line, $\Delta\sim 2\sqrt{m/M_N TE}$, at room temperature, $\Delta\sim 0.03\sqrt{E}$, where $E$ is in electron volts. At $\Delta>\Gamma\sim 0.03-0.1$ eV the effect is suppressed by a factor of about $\Delta/\Gamma$.

Let us examine the rotation of the spin of the neutron around its momentum direction. We assume that the neutron is moving along the $z$ axis and that at $z=0$ the spin is directed along the $x$ axis. After a distance $l$ is traversed, the components of the spinor acquire different phase shifts:

$$\frac{1}{\sqrt{2}}\begin{pmatrix}1\\1\end{pmatrix}\to\frac{1}{\sqrt{2}}\begin{pmatrix}\exp(ikn_+l)\\\exp(ikn_-l)\end{pmatrix}=\frac{1}{\sqrt{2}}\exp(ikn_+l)\begin{pmatrix}1\\e^{-i\psi}\end{pmatrix},\qquad \psi=kl\,\mathrm{Re}\,(n_+-n_-)=-\frac{2\pi Ng\Gamma_{\mathrm{n}}^{(p)}}{k_p^2}Pl\frac{E-E_p}{(E-E_p)^2+(\Gamma^2/4)}.\tag{4.8}$$

The resulting spinor corresponds to a spin which has rotated by an angle $-\psi$ around the $z$ axis. It is also simple to see that an unpolarized beam acquires a longitudinal polarization:

$$a=-kl\,\mathrm{Im}\,(n_+-n_-)=-\frac{2\pi Ng\Gamma_{\mathrm{n}}^{(p)}}{k_p^2}Pl\frac{\Gamma/2}{(E-E_p)^2+\Gamma^2/4}.\tag{4.9}$$

Although $\psi$ and $a$ are proportional to the distance traversed, $l$, it is clear that $l$ cannot be much larger than the mean free path $l_0=[k\,\mathrm{Im}(n_++n_-)]^{-1}\sim 1-5$ cm.

Suitable nuclei for such experiments are those which on the one hand have an isolated low-lying $p$-wave resonance but on the other hand a rather closely spaced spectrum of compound resonances. Examples are $^{113}$In, $^{117}$Sn, $^{119}$Sn, $^{139}$La, $^{232}$Th, and $^{238}$U. Let us examine the order of magnitude expected for these effects near $p$-wave resonances.

1) The relative difference between the capture cross sections for right-hand- and left-hand-polarized neutrons is

$$P=\frac{\sigma_+^a-\sigma_-^a}{\sigma_+^a+\sigma_-^a}\sim 10^{-2}.\tag{4.10}$$

2) The rotation of the neutron spin is

$$\psi\left(E_p-\frac{\Gamma}{2}\right)-\psi\left(E_p+\frac{\Gamma}{2}\right)\sim(10^{-2}-10^{-3})\frac{l}{l_0}\ \text{rad}.\tag{4.11}$$

3) The longitudinal polarization of the neutrons (or the relative difference between the transmission probabilities for right-hand- and left-hand-polarized neutrons) is

$$a(E_p)\sim(10^{-2}-10^{-3})\frac{l}{l_0}.\tag{4.12}$$

The values of $\psi$ and $a$ are several times smaller than $P$, because of the elastic scattering of neutrons ($\sigma_0\sim 5-10$ b), which reduces the mean free path. For the resonances under discussion here, we would have $\sigma^a/(\sigma^a+\sigma_0)\sim 0.2-0.5$.

These resonance effects are sharply dependent on the neutron energy. The scale dimension along the energy scale is $\sim\Gamma\sim 0.03-0.1$ eV, and this energy determines the permissible energy spread of the neutron beam. For a beam with an energy spread, $\psi$ and $a$ are inversely proportional to the spread.

We wish to emphasize that there are two factors which are responsible for the large magnitude of these effects. First, there is the kinematic enhancement, which occurs because the $s$-wave admixture amplitude is larger by a factor of $1/kR$ than the primary $p$-wave amplitude. Second, there is the dynamic enhancement of $P$-odd mixing in a compound nucleus.

### b) Magnitude of the effects for thermal neutrons. Comparison with experiment

We note at the outset that the $p$-wave and $s$-wave resonance generally make comparable contributions to the spin rotation angle in the thermal region, $\psi_t$. If there is an $s$-wave resonance close to the thermal region, however, the mean free path is small, reducing the observable effect. We will therefore consider the case in which there is a single $p$-wave resonance near the thermal region.

Working from (4.8) and (4.9), we can easily find an estimate of the effect in the thermal region ($E=0$). With $\Gamma\sim 0.1-0.03$ and $E_p\sim 1-10$ eV we find $\psi_t\sim 10^{-4}-10^{-5}l/l_0$ rad and $a_t\sim 10^{-6}-10^{-8}l/l_0$.

The first measurements of $\psi_t$ and $a_t$ in tin have recently been published by Forte *et al.*[13] They were seeking an effect of a single-particle $p$-wave resonance in $^{124}$Sn. As a control, they also carried out measurements in a natural mixture of tin isotopes. As it turned out, however, they failed to find an effect[5)] in $^{124}$Sn,

$$\varphi_t\,(^{124}\mathrm{Sn})=(0.48\pm 1.49)\cdot 10^{-6}\ \mathrm{rad/cm},\tag{4.13}$$

while for the control sample the angle turned out to be nonzero:

$$\varphi_t\,(\text{natural Sn})=(4.95\pm 0.93)\cdot 10^{-6}\ \mathrm{rad/cm}.\tag{4.14}$$

A detailed study revealed that the effect was attributable to $^{117}$Sn:

$$\varphi_t(^{117}\mathrm{Sn})=(36.7\pm 2.7)\cdot 10^{-6}\ \mathrm{rad/cm},\tag{4.15}$$

where there is a $p$-wave compound resonance at an energy of 1.32 eV. The measured value of $\varphi_t(^{117}\mathrm{Sn})$ agrees well with the above estimates.[40] Stodolsky[41] has also suggested that the large magnitude of the effect in $^{117}$Sn might be attributed to the proximity to the thermal region of a $p$-wave compound resonance.

It should be noted that an effect slightly smaller than that in $^{117}$Sn can be expected in $^{119}$Sn (provided, of course, that its spin is $J<2$), where there is also a $p$-wave resonance at 6.2 eV. We do not rule out the possibility that this isotope is responsible for the "ex-

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[5)] In (4.13)–(4.15) we have used the notation $\varphi_t$ instead of $\psi_t$ because of an ambiguity in the definition of the $\varphi_t$ (as discussed further on in the text proper). We wish to thank V. E. Bunakov and V. P. Gudkov for a discussion of this question of the sign of $\varphi_t$.

cess" effect observed in the natural mixture of isotopes.

Forte *et al.*[13] also measured the quantity

$$a_t\,(^{117}\mathrm{Sn}) = (-1.63 \pm 0.67)\cdot 10^{-6}\ \mathrm{cm}^{-1},$$
$$a_t\,(\text{null test}) = (-0.50 \pm 0.89)\cdot 10^{-6}\ \mathrm{cm}^{-1}. \qquad (4.16)$$

The relationship between the signs of $\psi_t$ and $a_t$ tells us which resonance (one with a positive or negative energy) is primarily responsible for the effect ($\psi_t/a_t = -2E_p/\Gamma$). Unfortunately, there is some ambiguity regarding the determination of the sign of $\varphi_t$ in the experimental paper by Forte *et al.*[13] [in Eq. (3) of that paper, the angle $\varphi$ was defined with a sign different from that in the text immediately preceding (3)]. If the definition of $\varphi_t$ differs from that of $\psi_t$ in sign, then the effect is due primarily to a resonance which lies below the neutron threshold. Working from the average value of the total width of the above-threshold resonance ($\Gamma = 0.08$ eV) and from the ratio $\psi_t/a_t$, we find $E_p \approx -1$ eV. The existence of a second $p$-wave resonance so close to the thermal point (the first is at $E_p = 1.32$ eV) seems unlikely (although, of course, we cannot rule out this possibility), since the typical distance between resonance in $^{117}$Sn is more than 20 eV. Under the circumstances we will assume that it is nevertheless the 1.32-eV resonance which makes the major contribution to the effect; i.e., we assume $\psi_t(^{117}\mathrm{Sn}) = \varphi_t$.

Working from the value of $\psi_t$ we can generate more accurate estimates of these effects in the vicinity of the resonance. For the resonance $E_p = 1.32$ eV, the neutron width is $g\Gamma_n^{(p)} = (1 \pm 0.5)\cdot 10^{-7}$ eV, according to Mughabghab and Garber.[42] From (4.4) and (4.8) we then find

$$P = \frac{\sigma_+^a - \sigma_-^a}{\sigma_+^a + \sigma_-^a} = (1.3 \pm 0.7)\cdot 10^{-2}. \qquad (4.17)$$

Unfortunately, the total width $\Gamma$ is not known for the 1.32-eV resonance. On the basis of the widths of higher resonances,[42] we will set $\Gamma = 0.08$ eV. The elastic cross section in $^{117}$Sn is $\sigma_0 \approx 5$ b. The total cross sections $\sigma_{\mathrm{tot}}$ in the thermal region and in the $p$-wave resonance turn out to be approximately equal at 7–8 b. Hence the mean free path is $l_0 \approx 4$ cm. Finally, for the quantities in which we are interested we find the values

$$\left.\begin{aligned} a_t = \psi_t \frac{\Gamma}{2E_p} &= -1.1\cdot 10^{-6}\ \mathrm{cm}^{-1} \ \text{ or } \ -4\cdot 10^{-6}\frac{l}{l_0},\\ \psi\left(E_p - \frac{\Gamma}{2}\right) - \psi\left(E_p + \frac{\Gamma}{2}\right) &= 1.2\cdot 10^{-3}\ \mathrm{rad/cm} \ \text{ or } \ 5\cdot 10^{-3}\frac{l}{l_0}\ \mathrm{rad},\\ a(E_p) &= -1.2\cdot 10^{-3}\ \mathrm{cm}^{-1} \ \text{ or } \ -5\cdot 10^{-3}\frac{l}{l_0}. \end{aligned}\right\} \qquad (4.18)$$

In these calculations we have used only the single assumption that $\psi_t$ is dominated by the 1.32-eV resonance. If we adopt the further assumption that the $s$-wave resonance nearest the thermal region has the same angular momentum $J$ as the 1.32-eV $p$-wave resonance, then we can determine the extent to which these states are mixed. In $^{117}$Sn, the cross section for the $(n,\gamma)$ reaction for thermal neutrons is 2.6 b. It is not difficult to see that known positive-energy resonances contribute $\lesssim 0.2$ b to this cross section. The cross section at the thermal point is therefore determined by a resonance with a negative energy. Working from the average widths of the resonances, we find

$$E_s \approx -10\ \mathrm{eV},\quad |\langle s|H_W|p\rangle| \approx 0.5\cdot 10^{-3}\ \mathrm{eV},$$
$$|\varepsilon| = |\alpha| = \left|\frac{\langle s|H_W|p\rangle}{E_s}\right| \approx 0.5\cdot 10^{-4} \simeq 2\cdot 10^2 G m_\pi^2. \qquad (4.19)$$

## 5. ENHANCEMENT OF PARITY-BREAKING EFFECTS IN THE (n, γ) REACTION

We will begin with an effect which was actually already discussed in the preceding section: the difference between the cross sections for the capture of right-hand- and left-hand-polarized neutrons into a $p$-wave resonance. It would apparently be convenient to measure the difference in the $(n,\gamma)$ reaction from the difference in the $\gamma$ count rates, since there would be no suppression by elastic scattering here. According to (4.10) the relative magnitude of the effect is ~$10^{-2}$.

We now consider the circular polarization of $\gamma$ rays in the $(n,\gamma)$ reaction in the capture of unpolarized neutrons into a $p$-wave compound resonance. Because of the difference between the cross sections $\sigma_+^a$ and $\sigma_-^a$, the intermediate compound nucleus is longitudinally polarized. Upon decay, this polarization is transferred to the $\gamma$ ray. We are thus dealing with an $(\mathbf{s}_\gamma\mathbf{p}_\gamma)(\mathbf{p}_\gamma\mathbf{p}_n)$ correlation; i.e., the degree of circular polarization is $\bar{P}_\gamma \propto \cos\theta$, where $\theta$ is the angle between the momenta of the $\gamma$ ray and the neutron. In order of magnitude, the circular polarization is $\bar{P}_\gamma \sim P \sim 10^{-2}$. If we assume, for example, that the 1.32-eV resonance in $^{117}$Sn has $J^P = 1^-$, then for a transition from this resonance to the $J^P = 0^+$ ground state of $^{118}$Sn we have

$$P_\gamma \sim P\cos\theta = (1.3 \pm 0.7)\cdot 10^{-2}\cos\theta. \qquad (5.1)$$

We turn now to the "classical" parity-breaking effects in the $(n,\gamma)$ reaction: the angular asymmetry in the $\gamma$ emission direction [the $(\mathbf{s}_n\mathbf{p}_\gamma)$ correlation] and the circular polarization of the $\gamma$ rays [the $(\mathbf{s}_\gamma\mathbf{p}_\gamma)$ correlation]. These correlations may also be enhanced near a $p$-wave resonance. To avoid writing out the lengthy equations for the arbitrary case, we consider a specific example for which experimental data are available[43,44]: the reaction $^{117}\mathrm{Sn}(n,\gamma)$, with a transition to the $0^+$ ground state of $^{118}$Sn. The $\gamma$ transition can come from the $1^+$ and $1^-$ compound levels. When the weak interaction is taken into account, the overall reaction amplitude can be written

$$\frac{\langle 0^+|\mathrm{M1}|1^+\rangle\langle 1^+|T_s|\mathrm{n}\rangle}{E - E_+ + (i\Gamma_+/2)} + \frac{\langle 0^+|\mathrm{E1}|1^-\rangle\langle 1^-|T_p|\mathrm{n}\rangle}{E - E_- + (i\Gamma_-/2)}$$
$$+ \frac{\langle 0^+|\mathrm{E1}|1^-\rangle\langle 1^-|H_W|1^+\rangle\langle 1^+|T_s|\mathrm{n}\rangle}{[E - E_- + (i\Gamma_-/2)][E - E_+ + (i\Gamma_+/2)]}$$
$$+ \frac{\langle 0^+|\mathrm{M1}|1^+\rangle\langle 1^+|H_W|1^-\rangle\langle 1^-|T_p|\mathrm{n}\rangle}{[E - E_- + (i\Gamma_-/2)][E - E_+ + (i\Gamma_+/2)]}. \qquad (5.2)$$

The fourth term in this equation is always smaller than the third and thus inconsequential. The interference of the second term with the third leads to the correlation $(\mathbf{s}_\gamma\mathbf{p}_\gamma)(\mathbf{p}_\gamma\mathbf{p}_n)$, which we have already discussed. An interference of the first term with the third is responsible for the $P$-odd correlations $(\mathbf{s}_\gamma\mathbf{p}_\gamma)$ and $(\mathbf{s}_n\mathbf{p}_\gamma)$. Unless we are very close to the $p$-wave resonance, we may ignore the second term. For the degree of circular polarization and the asymmetry parameter in the $\gamma$ angular distribution $W(\theta) = 1 + a\cos\theta$ we find

$$P_\gamma = a = 2\,\mathrm{Re}\left(\frac{\langle 1^-|H_W|1^+\rangle}{E - E_-}\,\frac{\langle 0^+|\mathrm{E1}|1^-\rangle}{\langle 0^+|\mathrm{M1}|1^+\rangle}\right). \qquad (5.3)$$

We wish to call attention to the fact that the energy denominator in this expression is of the form $E - E_-$, rather than $E_+ - E_-$. Therefore, under the condition $|E - E_-| \ll |E_+ - E_-|$, there will be an additional resonant enhancement by a factor $(E_+ - E_-)/(E - E_-)$ in comparison with the usual estimate for the mixing ratio.

An asymmetry in the angular distribution of $\gamma$ rays in the reaction $^{117}Sn(n,\gamma)^{118}Sn$ for thermal neutrons has been found experimentally: $a = (8.9 \pm 1.5) \cdot 10^{-4}$ according to Ref. 43 and $(4.4 \pm 0.6) \cdot 10^{-4}$ according to Ref. 44. Let us attempt to compare this value with the results from neutron optics. Unfortunately, the spin of the 1.32-eV resonance is not known. We have essentially eliminated the possibility $J = 2$ already, when we assumed that this resonance contributed to the rotation of the neutron spin. If $J = 0$, then this resonance does not contribute to the angular asymmetry of the $\gamma$ rays, and there is no direct relationship between the effects. All we can do is use the value of $\varepsilon$ from (4.19) and write the standard estimate $|a| \sim |2\varepsilon \mathrm{E1}/\mathrm{M1}| \approx 10^{-4}\ |\mathrm{E1}/\mathrm{M1}|$. If $J = 1$, on the other hand, then the effects are related. Using (5.3) and (4.19), we find

$$|a| \approx \left| 2\varepsilon \frac{E_s}{E_p} \frac{E1}{M1} \right| \approx 8 \cdot 10^{-4} \left| \frac{E1}{M1} \right|. \qquad (5.4)$$

In the derivation of this expression it was implicitly assumed that the amplitudes for electromagnetic transitions to the $^{118}Sn$ ground state are determined at the thermal point by the nearest $s$-wave and the nearest $p$-wave resonances ($|s_0\rangle = |1^+, E_s \approx -10\ \mathrm{eV}\rangle$, $|p_0\rangle = |1^-, E_p = 1.32\ \mathrm{eV}\rangle$). Taking all these resonances into account, we would write

$$a = -2\,\mathrm{Re}\,\frac{\sum\limits_{s,p}\left(\frac{\langle 0^+|E1|p\rangle\langle p|H_W|s\rangle\langle s|T_s|n\rangle}{E_s E_p}\right)}{\sum\limits_{s}\left(\frac{\langle 0^+|M1|s\rangle\langle s|T_s|n\rangle}{E_s}\right)}. \qquad (5.5)$$

We could cite arguments according to the sum which in the numerator is actually saturated by the nearest resonances. For the denominator, on the other hand, a large number of terms are important (up to $|E_s| \sim 1$ MeV). Taking these circumstances into account, we easily see that we again find Eq. (5.4), but now the amplitude M1 $= \langle 0^+|M1|s_0\rangle$ in this equation should be replaced by a certain effective overall amplitude for the M1 transition at the thermal point:

$$\widetilde{M1} = \langle 0^+|M1|s_0\rangle + \sum_{s \neq s_0} \frac{E_{s0}}{E_s} \langle 0^+|M1|s\rangle \frac{\langle s|T_s|n\rangle}{\langle s_0|T_s|n\rangle}. \qquad (5.6)$$

According to rough estimates, the sum in this expression is of the same order of magnitude as the first term, or perhaps even larger by a factor of two or three.[6)] For the same reason, the ratio $|\mathrm{E1}/\widetilde{\mathrm{M1}}| \sim 1$ which follows from a comparison of (5.4) with the experimental asymmetry seems quite reasonable.

To avoid any possible confusion, we stress the fact that the comparatively large contribution of the remote resonances to the M1 amplitude is a consequence of our consideration of the $\gamma$ transition to the $^{118}Sn$ ground state. The total cross section for the reaction $^{117}Sn(n,\gamma)$, in contrast, is determined primarily by the nearest $s$-wave resonance.

We wish to thank I. B. Khriplovich for a discussion which extended to essentially all the questions covered in this review and which we found exceedingly important.

## APPENDIX. PERTURBATION-THEORY DESCRIPTION OF THE NUCLEUS FROM THE COMPOUND STATE TO THE COLD STEP

For a better understanding of how the part of the compound-nucleus wave function corresponding to the cold step arises [see (2.2) and (2.3)], it is useful to consider a simple but quite realistic model: fission in perturbation theory. This model will also cast some light on the origin of the phase shifts in the amplitude *for the transition to the cold step* [*see* (2.11)].

Upon capture into a compound state, a neutron excites some of the nucleons of the nucleus. The excited nucleons and holes move in a certain average core potential and interact strongly with each other. Their interaction with the core vibrations, on the other hand, will be assumed small here and will be dealt with by perturbation theory. From the formal standpoint, this problem is analogous to that of the interaction of a gas of quasiparticles with a heavy particle which is moving in a potential well. Fission corresponds to the emission of the particle from the well. We assume that fission occurs just above the barrier; i.e., fission occurs only if the gas "cools off" completely and transfers energy to the core vibrations. Figure 1 shows a diagram corresponding to this process. We are interested in the lowest order of perturbation theory in terms of the interaction of the quasiparticles with the core vibrations. Since there are no excited nucleons in the final state, a nucleon must annihilate with a hole each time there is an interaction with vibrations. In other words, if $n$ nucleons are initially excited, then fission occurs in perturbation theory of order $n$. The amplitude corresponding to Fig. 1(a) is

$$A_n(E) = \sum_{\alpha_i k_i}{}' \frac{\langle a|V|k_1\alpha_1\rangle}{E - E_1 + (i\Gamma_1/2)} \times \cdots \frac{\langle k_{n-1}, \alpha_{n-1}|V|0, n\rangle}{E - E_{n-1} + (i\Gamma_{n-1}/2)}. \qquad (A1)$$

Here $|0,n\rangle = |0\rangle|n\rangle$ is the initial state; $|0\rangle$ is the core wave function; $|n\rangle$ represents a gas of $n$ nucleons and $n$ holes; $|k_i, \alpha_i\rangle \equiv |k_i\rangle|\alpha_i\rangle$, $|\alpha_i\rangle$ is the wave function of $i$

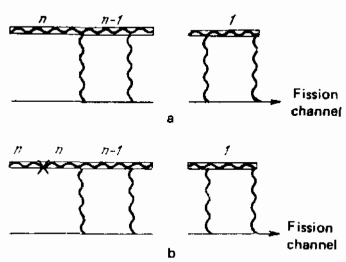


FIG. 1. Heavy line—Gas of nucleons and holes; solid line—core; wavy line—interaction of a nucleon with core vibrations; cross—weak interaction.

[6)] By way of comparison we note that in the reaction $^{113}Cd(n,\gamma)^{114}Cd(0^+)$ the sum is apparently much smaller than the first term.

nucleons and $i$ holes; $|k_i\rangle$ is the core state; $|a\rangle$ is the final state, a fixed fission channel; $E$ is the initial energy; and $E_i$ and $\Gamma_i$ are the energies and widths of the intermediate states. The sum runs over all the intermediate states, of which there are many. It is not difficult to see, however, that the predominant terms are those which correspond to the minimum possible "energy nonconservation" in the intermediate states, since the transitions which are dynamically enhanced are those which have the smallest energy denominators in Eq. (A1).

The parity-breaking weak interaction can be incorporated in any step of the process in Fig. 1(a), but a weak interaction in the initial step will obviously make the predominant contribution [Fig. 1(b)]. In the initial step, the number of excited nucleons is at a maximum, so that the level spectrum is at its densest, and the dynamic-enhancement factor is at its highest. The admixture amplitude corresponding to Fig. 1(b) is

$$C(E)=\sum_{\nu} A_{\bar{\eta}}(\nu, E)\frac{\langle 0, \nu|H_W|0, n\rangle}{E-E_\nu+(i\Gamma_\nu/2)}, \qquad (A2)$$

where $H_W$ is the operator representing the parity-breaking weak interaction, and $A_{\bar{\eta}}(\nu,E)$ is the amplitude for fission from the admixture level; this amplitude is of the same form as $A_\eta(E)$ in Eq. (A1). The sum runs over all the levels of opposite parity of the system of $n$ nucleons. Only the nearest levels, however, make a dynamically enhanced contribution. Since the admixture amplitude has the same structure as the primary amplitude, it is clear that there is no further suppression of any sort of the admixture amplitude because of the transition from the hot step to the cold step. We thus again find Eqs. (2.5) and (2.7). In perturbation-theory terms, the phase shifts in the transition amplitudes [$A\sim\sqrt{\Gamma}e^{i\varphi}$; see (2.11)] arise when the widths of the intermediate states are incorporated in the diagrams in Fig. 1. Since the spectrum is less dense in these states than in the initial compound nucleus, the energy dependence of the phase shifts is smooth in comparison with the dependence $[E-E_\nu+(i\Gamma_\nu/2)]^{-1}$ associated with the initial compound nucleus.

In connection with the discussion of the $T$ invariance in Subsec. 2b, we note that it is completely obvious in perturbation theory that if all the $\Gamma$ and $\Gamma_\nu$ vanish then we have $\cos(\varphi-\varphi_\nu)=\pm 1$.


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