Testing T-Odd, P-Even Interactions with γ Rays from Neutron p-Wave Resonances

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A new method for the study of time reversal violation is described. It consists of measurements of the forward-backward asymmetry in individual gamma-ray transitions resulting from unpolarized neutron capture in p-wave resonances. An experiment with a ^{113}Cd target performed at the Dubna pulsed neutron source has been analyzed and a limit on the time reversal odd, parity even interaction extracted. The possibility of experiments using the powerful pulsed neutron source at Los Alamos is considered.

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CP nonconservation or T (time reversal) noninvariance was observed in the decay of neutral kaons long ago [1]. Attempts to search for this phenomenon in other processes have failed. Recently some possible tests using neutron optics experiments were discussed in a paper by Weidenmuller [2] and references therein. Interest in the neutron-nucleus interaction was stimulated by the discovery [3,4] of large P- (parity-) nonconserving effects near neutron p-wave resonances. These effects appeared to be enhanced by 3 to 6 orders of magnitude compared to the single-particle estimate $Gm_\pi^2 \sim 10^{-7}$ (in units of $\hbar = c = 1$ and $G = 10^{-5} m_p^{-2}$). The enhancement was explained in terms of the P-nonconserving mixing of compound nuclear states with opposite parity (mixing of the p-wave resonance with the nearest s-wave resonances [5,6]). The concept of the spreading width of the weak interaction was brought into practice based on recent experimental results [7].

It was shown by Bunakov and Gudkov [8,9] that possible T-noninvariant effects may be enhanced near p-wave resonances in the same manner as P-nonconserving effects. It has been suggested that a serach be made for P-conserving, T-noninvariant effects in polarized neutron transmission through aligned nuclear targets [10–12]. Such an effect would occur due to an $\mathbf{s}\cdot[\mathbf{k}\times\mathbf{I}](\mathbf{k}\cdot\mathbf{I})$ term in the neutron-nucleus elastic forward scattering amplitude. Here $\mathbf{s}$ and $\mathbf{I}$ are the neutron and nucleus spins and $\mathbf{k}$ is neutron momentum. Bunakov [9] showed that the maximum value of the experimental polarization asymmetry in the cross section β is connected to the T-noninvariant matrix element v_p^T and to the p-wave resonance level spacing D_p by the approximate relation

$$\beta \approx v_p^T / D_p \, . \quad (1)$$

Here $iv_p^T = \langle H^T \rangle_{12} = -\langle H^T \rangle_{21}$ is the purely imaginary matrix element of the T-noninvariant, P-conserving interaction between compound nuclear wave functions of the considered p-wave resonance (labeled 1) and its nearest-neighboring p-wave resonance (labeled 2). Estimates given by Bunakov [9] indicate the enhancement factor may be as big as 10^3. The estimate is that $v_p^T/D_p \sim 10^3 \phi$, where ϕ is roughly the strength of the T-noninvariant, P-conserving nuclear interaction relative to the T-invariant one. The new theoretical limits on ϕ are very low [13]. The experimental upper limit on ϕ is between 10^{-3} and 10^{-4}, e.g., as given in Ref. [9] or in the recent aligned target experiment [14]. Therefore, in the framework of Bunakov's hypothesis, any experimental measurement of β near a p-wave resonance with accuracy better than 10^{-1} may decrease this limit.

It has been noted [15] that the value of v_p^T/D_p in p-wave resonances may be obtained from a quite different kind of experiment, namely, from the measurement of the forward-backward asymmetry in the yield of gamma rays from individual transitions in unpolarized neutron capture reactions. Such measurements have already been made [16–18] for p-wave resonances in targets of ^{113}Cd and ^{117}Sn. Here we present the results of our analysis of the ^{113}Cd capture experiment after giving the details of our approach to the analysis which is based on the hypothesis of Ref. [9].

In our approach to the calculation of experimental observables we will use an S-matrix scattering formalism. We wish to start from the expression

$$\delta S_J = S_J(1\tfrac{1}{2} \to 1\tfrac{3}{2}) - S_J(1\tfrac{3}{2} \to 1\tfrac{1}{2}) \, ,$$

where $S_J(lj, l'j')$ is the scattering matrix element which corresponds to the transition of p-wave neutrons ($l = l' = 1$) in a resonance with spin J from a channel with neutron total angular momentum $j = \frac{1}{2}$ to a channel with $j' = \frac{3}{2}$ and vice versa. Using the explicit forms for the scattering matrix elements of Ref. [19] we get

$$\delta S_J = \frac{2\,\mathrm{Im}[g_n(1\tfrac{1}{2}) g_n^*(1\tfrac{3}{2})]}{E - E_p + i\Gamma_p/2} \, . \quad (2)$$

Here E_p and Γ_p are the energy and the total width of the p-wave resonance and $g_n(lj)$ is the neutron partial width amplitude $[\Gamma_n(lj) = |g_n(lj)|^2]$. By definition

$$g_n = \alpha \langle \varphi_{JMlj} | \Psi_{JM} \rangle \, , \quad (3)$$

where Ψ_{JM} is the compound nuclear wave function, the functions φ_{JMlj} describe the neutron and nucleus in the entrance and exit channels, and the coefficient α depends on the normalization conditions. We neglect potential scattering phases which are small for low-energy resonance neutrons.

If T invariance holds, the amplitude $g_n(lj)$ is real [19]. A T-noninvariant, P-conserving interaction H_T gives phases $\delta_n(lj)$ to these amplitudes and mixes the compound nuclear wave functions. Using the hypothesis of Bunakov [9], we have for an isolated p-wave resonance

$$\Psi_{JM}=\Psi_{JM}^{(1)}+\frac{\langle H^T\rangle_{21}}{E-E_{p2}+i\Gamma_{p2}/2}\Psi_{JM}^{(2)} \tag{4}$$

and

$$g_n(1j)=g_n^{(1)}\exp[i\delta_n^{(1)}(1j)]-i\frac{v_p^T}{E-E_{p2}+i\Gamma_{p2}/2}g_n^{(2)}\exp[i\delta_n^{(2)}(1j)]\,, \tag{5}$$

where $g_n^{(1,2)}\equiv g_n(1j)^{(1,2)}$ and the superscripted indexes 1 and 2 correspond to the considered p-wave resonance and to its nearest-neighboring p-wave resonance, respectively. Taking into account that $|D_p|=|E_{p1}-E_{p2}|\gg\Gamma_{p2}$ we obtain in first order in H_T near resonance 1,

$$\delta S_J=-2\alpha\frac{g_n^{(1)}(1\tfrac{1}{2})\mathrm{Im}\langle\varphi_{JM1(3/2)}|\Psi_{JM}^{(1)}\rangle-g_n^{(1)}(1\tfrac{3}{2})\mathrm{Im}\langle\varphi_{JM1(1/2)}|\Psi_{JM}^{(1)}\rangle}{E-E_{p1}+i\Gamma_{p1}/2}+\frac{2v_p^T}{D_p}\frac{g_n^{(1)}(1\tfrac{1}{2})g_n^{(2)}(1\tfrac{3}{2})-g_n^{(1)}(1\tfrac{3}{2})g_n^{(2)}(1\tfrac{1}{2})}{E-E_{p1}+i\Gamma_{p1}/2} \tag{6}$$

in accordance with Eq. (6) of Ref. [9]. Since $\alpha\,\mathrm{Im}\langle\varphi_{JMlj}|\Psi_{JM}\rangle\cong\phi g_n(lj)$ and $v_p^T/D_p\cong 10^3\phi$, the second term in Eq. (6) dominates. This leads to Eq. (1) for the cross-section asymmetry β in an isolated p-wave resonance (see also Ref. [20]) in a transmission experiment with polarized neutrons.

Now we consider the capture of neutrons by a spin $I=\frac{1}{2}$ nucleus, ^{113}Cd, for example, with the individual gamma rays being detected. The asymmetry in gamma-ray emission in $1^+,1^-\rightarrow 0^+$ transitions to the ground state of the product nucleus ^{114}Cd is caused by the interference of $M1$- and $E1$-electromagnetic waves. This can be described by the second term in the following expression for the differential cross section:

$$d\sigma(\mathbf{n}_\gamma,E)/d\Omega=A_0(E)+A_1(E)(\mathbf{n}_\gamma\cdot\mathbf{n}_k)+A_2(E)P_2(\mathbf{n}_\gamma\cdot\mathbf{n}_k)\,, \tag{7}$$

where $\mathbf{n}_\gamma$ and $\mathbf{n}_k$ are unit vectors in the direction of the photon and neutron momenta and $P_2(\mathbf{n}_\gamma\cdot\mathbf{n}_k)$ is the second Legendre polynomial. In our case the $A_2(E)$ term does not contribute substantially and will be dropped from further discussion. The coefficients $A_0(E)$ and $A_1(E)$ depend on the neutron energy and have the form

$$A_0(E)=(g_J/4k^2)\left(|f_s(E)|^2+\sum_j|f_{pj}(E)|^2\right), \tag{8}$$

$$A_1(E)=-(g_J/2k^2)\mathrm{Re}\{f_s^*(E)[f_{p(1/2)}(E)-f_{p(3/2)}(E)/\sqrt{2}]\}\,, \tag{9}$$

where

$$f_s(E)=\frac{g_n^s(0\tfrac{1}{2})g_\gamma^s}{E-E_s+i\Gamma_s/2},\quad f_{pj}(E)=\frac{g_n^p(1j)g_\gamma^p}{E-E_p+i\Gamma_p/2} \tag{10}$$

are the amplitudes of the (n,γ) reaction proceeding via s- and p-wave neutron resonances, and $g_{J=1}=\frac{3}{4}$ is the spin statistical factor. Here the radiative amplitudes g_γ^s and g_γ^p of electromagnetic $M1$ and $E1$ transitions are introduced in the same manner described in Ref. [21]:

$$g_\gamma^{s(p)}\sim\int d^3r\langle\Psi_{JM}^{s(p)}|\mathbf{A}_{1M}(\mathbf{r},m(e))\hat{\mathbf{j}}(\mathbf{r})|0^+\rangle\,, \tag{11}$$

where $\mathbf{A}_{JM}(\mathbf{r},m(e))$ are multipole fields and $\hat{\mathbf{j}}(\mathbf{r})$ is the current operator. If T invariance holds, these radiative amplitudes are real. As a consequence of T-noninvariant mixing of the p-wave resonance wave functions [Eq. (4)] we obtain

$$g_\gamma^p=g_\gamma^{p1}\exp(i\delta_\gamma^{p1})+i\frac{v_p^T}{E-E_{p2}-i\Gamma_{p2}/2}g_\gamma^{p2}\exp(i\delta_\gamma^{p2})\,. \tag{12}$$

The interaction H^T may mix s-wave compound nuclear states as well, this time through the matrix element v_s^T. Therefore the amplitudes $g_n^s(0\frac{1}{2})$ and g_γ^s have to be obtained in forms similar to Eqs. (5) and (12). Taking into account the condition $|E_{s1}-E_{p1}|\gg\Gamma_{s1}$ and keeping among the T-noninvariant terms only those linearly dependent on v_p^T/D_p and v_s^T/D_s we will have

$$A_0(E)=\left(\frac{g_J}{4k^2}\right)\left(\frac{\Gamma_n^{s1}(0\tfrac{1}{2})\Gamma_\gamma^s}{(E-E_{s1})^2}+\frac{[\Gamma_n^{p1}(1\tfrac{1}{2})+\Gamma_n^{p1}(1\tfrac{3}{2})]\Gamma_\gamma^{p1}}{(E-E_{p1})^2+(\Gamma_{p1}/2)^2}\right) \tag{13}$$

and

$$A_1(E) = -\left(\frac{g_J}{2k^2}\right)\left(\frac{g_n^{s1}(0\tfrac{1}{2})[g_n^{p1}(1\tfrac{1}{2}) - g_n^{p1}(1\tfrac{3}{2})/\sqrt{2}]g_\gamma^{s1}g_\gamma^{p1}}{(E-E_{p1})^2 + (\Gamma_{p1}/2)^2}\right)\left(\frac{E - E_{p1} - \Delta E_{p1}}{E - E_{s1}}\right), \tag{14}$$

where

$$\Delta E_{p1} = \frac{\Gamma_{p1}}{2}\left[-\frac{\Gamma_{s1}/2}{E_{p1}-E_{s1}} - \frac{v_p^T}{D_p}\left(\frac{g_\gamma^{p2}}{g_\gamma^{p1}} - \frac{g_n^{p2}(1\tfrac{1}{2}) - g_n^{p2}(1\tfrac{3}{2})/\sqrt{2}}{g_n^{p1}(1\tfrac{1}{2}) - g_n^{p1}(1\tfrac{3}{2})/\sqrt{2}}\right) + \frac{v_s^T}{D_s}\left(\frac{g_\gamma^{s2}}{g_\gamma^{s1}} - \frac{g_n^{s2}(0\tfrac{1}{2})}{g_n^{s1}(0\tfrac{1}{2})}\right)\right]. \tag{15}$$

Equation (15) contains the primary result, namely, a T-noninvariant, P-conserving interaction which gives rise to a shift in the zero crossing of the fore-aft asymmetry in addition to the known first term of Eq. (15). This effect appears due to the polarization of the $E1$ and $M1$ quanta in the plane of the reaction. This explains the apparently surprising result that a T-violating shift is accessible to measurement with an unpolarized beam and target.

This experiment [17] has been performed using neutrons from the Dubna IBR-30 pulsed reactor. The reactor was operated in the booster mode as a multiplier of neutrons from the target of an electron linear accelerator. The duration of the electron pulse was 4.5 μs with pulses occurring at a 100 Hz repetition rate. A flight path of 52 m was used for the experiment. The flux on that flight path was equal to $10^4/E^{0.9}$ neutrons/cm^{2}eV s. The background was determined by the resonance absorber method using a tantalum sample for which the 4.3 and 10.3 eV resonances were "black." Two identical NaI(Tl) crystals (200 mm in diameter and 200 mm thick) were used as gamma-ray detectors. They were placed at 53° and 127° with respect to the beam direction at a distance of 40 cm from the ^{113}Cd sample. The sample was in the form of a 116 g, 70 mm diam cadmium disk 95% enriched in ^{113}Cd. The angular positions were chosen because the second Legendre polynomial, Eq. (7), vanishes totally at these angles. The energy resolution of the crystals permitted separation of the transitions to the ground state (9.04 MeV) and first excited state (8.48 MeV). The pulse height spectra were collected in 16 time gates ("windows") in the first run and then in 32 narrower gates during the second run. The typical energy width of these gates was 0.078 eV in the first run and 0.039 eV in the second run. The gate widths were small compared to the value of the total width $\Gamma = 0.16$ eV of p-wave resonance at $E_p = 7.0$ eV. For other experimental details see Ref. [18].

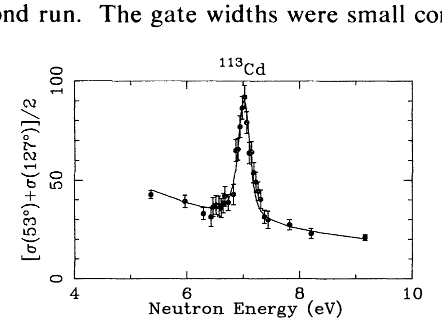

FIG. 1. The sum of gamma-ray yields in detectors placed at 53° and 127° with respect to the beam direction. Only gamma rays from an energy window corresponding to the direct transition to the ground state in ^{114}Cd were counted. The cross sections are given in arbitrary units.

After background subtraction, the number of counts in each detector (8.8 MeV threshold) was converted to a differential cross section. This conversion was performed for each neutron energy window. The resulting sum $[\sigma(53°)+\sigma(127°)]/2$ is plotted in Fig. 1 with arbitrary units. The error bars represent the statistical uncertainty associated with each data point. The solid line is the result of the analysis described below. The difference in the two cross sections, $[\sigma(53°)-\sigma(127°)]/2$, is shown in Fig. 2 in the same relative units. Both data sets were fitted using the CERN least-squares minimization program MINUIT and subprograms used to take into account the Doppler broadening of the resonance and the time-of-flight spectrometer resolution function.

To obtain the value of the resonance energy E_{p1} with high precision we fitted the data shown in Fig. 1 by Eq. (13) with the additional input of the known Doppler broadening and resolution function parameters. The values of Γ_{p1} and $\Gamma_n^{p1}(1\tfrac{1}{2})+\Gamma_n^{p1}(1\tfrac{3}{2})$ were obtained from Ref. [17]. Following the approach of Skoy and Sharapov [18] we replace all parameters of s-wave capture with 1, the value of the s-wave cross section at $E = E_{p1}$. An energy dependence of the form $\sigma_\gamma^s(E) = \sigma_\gamma^s(E_{p1})(E_{p1}/E)^{3/2}$ was used to calculate σ_γ^s at $E = E_{p1}$. This experimental energy dependence is weaker than that calculated using the parameters of the s-wave resonance at $E_s = 0.178$ eV. The measured total capture cross section near 7 eV was found to be at least a factor of 2 larger than the calculated value. Therefore $\sigma_\gamma^s(E_{p1})$ as well as the partial width of the 9.04 MeV transition, Γ_γ^{p1}, were treated as unknown parameters in addition to E_{p1}. The data in Fig. 2 were fitted by Eq. (14) modified to in-

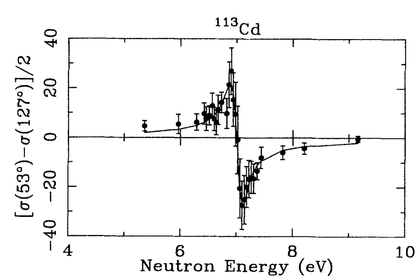

FIG. 2. The difference in cross section for the forward and backward angle detectors. The vertical axis has the same relative units as Fig. 1.

clude Doppler and resolution broadening. The free parameters in the fit were the energy shift ΔE_{p1} and the mixing ratio of the $g_n(1\frac{1}{2})$ and $g_n(1\frac{3}{2})$ amplitudes studied earlier by Alfimenkov *et al.* [17].

The result for the resonance energy and the resonance energy shift from the first run is $E_{p1}=7.0401\pm0.0048$ eV and $\Delta E_{p1}=-0.0005\pm0.0053$ eV with $\chi^2=4$. The large χ^2 was the reason for repeating the measurement. The second run (the data shown in Fig. 1 and Fig. 2) was performed with more regard to the background and with the narrow gates. The results of the second run give $E_{p1}=7.0412\pm0.0073$ eV and $\Delta E_{p1}=-0.0067\pm0.0113$ eV with $\chi^2=0.9$. The weighted value of the energy shift from both measurements is $\Delta E_{p1}=-0.0016\pm0.0062$ eV, consistent with time reversal invariance. The contribution $\Gamma_{p1}\Gamma_s/[4(E_{p1}-E_s)]$ of the nearest s-wave resonance ($E_s=0.178$ eV, $\Gamma_s=0.113$ eV) to ΔE_{p1} is 0.00066 eV and is therefore not considered. Neglecting the T-noninvariant interactions between s-wave levels and making the optimal choice for the neutron partial width amplitudes we may reduce Eq. (15) to the approximate relation $\Delta E_{p1}\approx(\Gamma_p/2)\beta$. Then the experimental upper bound for β is 0.08. The use of the dynamical enhancement factor in the relation $\beta\simeq10^3\phi$ sets an upper bound of 10^{-4} on ϕ. One must, however, consider that the enhancement factor given in Ref. [9] is statistically averaged. Some arguments for a lower value of the enhancement were given by Gudkov [22]. In a recent study [23] of the p-wave resonances in ^{113}Cd several close lying pairs of p-wave resonances were found. Pairs with small energy denominators should have large dynamical enhancements.

Experiments of this type may also be performed at the Los Alamos Neutron Scattering Center (LANSCE). The flux on the 60 m flight path at LANSCE is an order of magnitude larger than the flux on the 52 m flight path at the Dubna IBR-30 reactor. LANSCE has the further advantage that it is a spallation source using an 800 MeV proton beam in 250 ns bursts. The relatively narrow proton burst width at LANSCE may allow shorter flight paths to be used, with a corresponding increase in count rate. Alternatively, the narrow beam burst width allows one to study resonances up to approximately 500 eV. The higher intensity and better resolution of LANSCE will allow experiments of much higher precision to be performed and will allow the study of many resonances per nucleus.

We have presented a new technqiue for testing time reversal invariance using neutrons as a probe of the nuclear medium. The measurement of the zero crossing of the fore-aft asymmetry in individual gamma-ray transitions is a straightforward experimental task. The existence of dynamical enhancements in the compound nuclear system allows one to set relatively stringent limits on ϕ. Future experiments at the intense neutron source at Los Alamos should be able to improve the limit on β by an order of magnitude or more.

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