Parity mixed doublets in $A = 36$ nuclei

Mihai Horoi
*National Superconducting Cyclotron Laboratory and Department of Physics and Astronomy, Michigan State University, East Lansing, Michigan 48824
and Department of Theoretical Physics, Institute of Atomic Physics, P.O. Box MG - 6, Magurele, Bucharest, R-76900, Romania*
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The γ-circular polarizations (P_γ) and asymmetries (A_γ) of the parity forbidden $M1+E2$ γ decays, ^{36}Cl$^*(J^\pi = 2^-; T = 1; E_x = 1.95$ MeV$) \rightarrow {}^{36}$Cl$(J^\pi = 2^+; T = 1;$ g.s.) and ^{36}Ar$^*(J^\pi = 2^-; T = 0; E_x = 4.97$ MeV$) \rightarrow {}^{36}$Ar$^*(J^\pi = 2^+; T = 0; E_x = 1.97$ MeV), are investigated theoretically. We use the recently proposed Warburton-Becker-Brown shell-model interaction. For the weak forces we discuss comparatively different weak interaction models based on different assumptions for evaluating the weak meson-hadron coupling constants. The results determine a range of P_γ values from which we find the most probable values: $P_\gamma = 1.1 \times 10^{-4}$ for ^{36}Cl and $P_\gamma = 3.5 \times 10^{-4}$ for ^{36}Ar.

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Parity nonconservation (PNC) in the nucleon-nucleon interaction has been observed in the left-right asymmetry in $\vec{p}$-p scattering [1], in nucleon-nucleus scattering induced by polarized projectiles (such as $\vec{p}$ [2] or $\vec{n}$ [3]) in spontaneous α decay [4,5] and in the circular polarization [6–8] or asymmetry [9–12] (from polarized nuclei) of the radiation emitted in nuclear γ decay. There are also theoretical predictions for new PNC experiments in induced α decay [13,14] and asymmetry of the radiation emitted in nuclear γ decay [15,16]. The theoretical and experimental work in this field has been reviewed recently [11,12,17].

The controversy [11,12,18–20] in calculating weak meson-nucleon coupling constants in nuclei has greatly stimulated the investigation of possible experiments sensitive to different components of the PNC interaction Hamiltonian ($H_{\rm PNC}$) that depend linearly on seven such weak coupling constants: $h^{\Delta T}_{\rm meson}$: h^1_π, h^0_ρ, h^1_ρ, h^2_ρ, $h^1_{\rho'}$, h^0_ω, h^1_ω. Various linear combinations of these constants can, in principle, be extracted in different experiments, and among these are those for the parity mixed doublets (PMD) [11,16]. The most interesting PMD cases are those for which the PNC effect is enhanced due to a small energy difference between the two states and due to a favorable ratio of the transition probabilities. Since the PMD has definite isospins, the transition "filters out" specific isospin components of PNC weak interaction.

In the present paper a theoretical investigation of two new PMD cases in nuclei with $A = 36$ is presented. The first one, in ^{36}Cl, is given by the $J^\pi T = 2^-1$, E_x=1.951 MeV and $J^\pi T = 2^+1$, E_x=1.959 MeV levels (see Fig. 1). The second one, in ^{36}Ar, is given by the $J^\pi T = 2^-0$, E_x=4.974 MeV and $J^\pi T = 2^+0$, E_x=4.951 MeV lev-

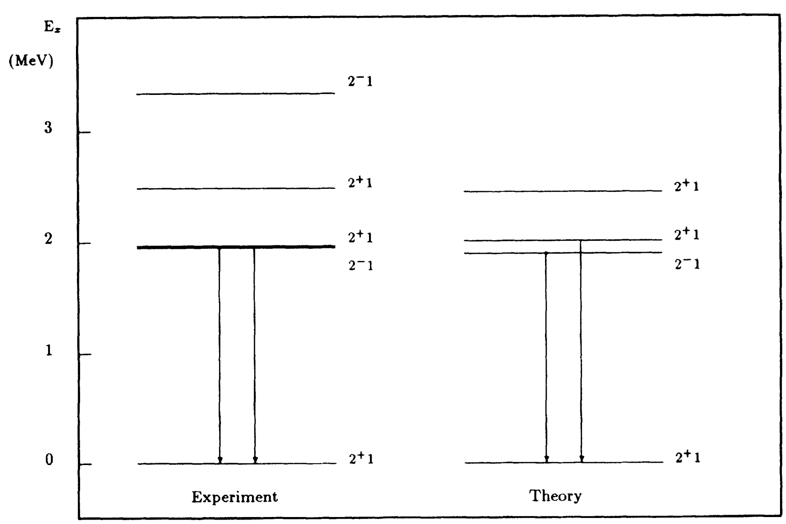

FIG. 1. Experimentally and theoretically calculated energies for the low $2^\pm1$ levels in ^{36}Cl. The first excited 2^+1 level has been artificially drawn 8 keV higher in order that the PMD be seen.

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els (see Fig. 2). In order to have an amplification of the PNC effect we look to the suppressed transitions from the $J^\pi T = 2^-1$ levels to the $J^\pi T = 2^+1$ ground state for ^{36}Cl and from the $J^\pi T = 2^-0$ level to the $J^\pi T = 2^+0$, E_x=1.97 MeV, level for ^{36}Ar. The corresponding PNC matrix elements were calculated with the shell-model code OXBASH, with the Warburton-Becker-Brown interaction [22] for $2s1d$-$2p1f$ model space. One of the cases considered here (^{36}Cl) has been investigated previously [15] with a much smaller valence model space which included only the $1d_{\frac{3}{2}}$ and $1f_{\frac{7}{2}}$ orbitals. Within this ($1d_{\frac{3}{2}}$, $1f_{\frac{7}{2}}$) small model space the one-body contribution to the PNC matrix element between the members of the doublet ($M_{\rm PNC}$) vanishes. Within the present model space the contribution of the one-body term dominates the theoretical $M_{\rm PNC}$. The goal of the present work is to calculate the PNC γ asymmetries and circular polarizations of the proposed gamma ray transitions, within different weak-interaction models, in order judge the experimental feasibility.

The degree of circular polarization of the emitted γ rays is given [see Ref. [23], Chap. 9, Sec. 3, Eq. (9.38)] by a sum of parity nonconserving (PNC) and parity conserving (PC) contributions:

$$P_\gamma(\cos\theta) \equiv \frac{W_{\rm right}(\theta) - W_{\rm left}(\theta)}{W_{\rm right}(\theta) + W_{\rm left}(\theta)} = (P_\gamma)_0 R_\gamma^{\rm PNC}(\cos\theta) + R_\gamma^{\rm PC}(\cos\theta), \tag{1}$$

where $R_\gamma^{\rm PC}$ is a parity conserving quantity discussed below,

$$(P_\gamma)_0 = 2\frac{M_{\rm PNC}}{\Delta E}\sqrt{\frac{b_+\tau_-}{b_-\tau_+}\left(\frac{E_\gamma^-}{E_\gamma^+}\right)^3}\sqrt{\frac{1+\delta_-^2}{1+\delta_+^2}} \tag{2}$$

and

$$R_\gamma^{\rm PNC}(\cos\theta) = \sqrt{\frac{1+\delta_-^2}{1+\delta_+^2}}\left(\sum_{\nu=0,2,4} P_\nu(\cos\theta)B_\nu(2)[F_\nu(1122) + F_\nu(2222)\delta_+\delta_+ + F_\nu(1222)(\delta_- + \delta_+)]\right)$$
$$\times\left(\sum_{\nu=0,2,4} P_\nu(\cos\theta)B_\nu(2)[F_\nu(1122) + F_\nu(2222)\delta_-^2 + 2F_\nu(1222)\delta_-]\right)^{-1}. \tag{3}$$

$R_\gamma^{\rm PNC}$ is a multiplier due to the existence of the orientation of the nuclear spin in the initial excited state when the mixing ratios do not vanish. In the above equations, δ_- is the $M2/E1$ mixing ratio, δ_+ is the $E2/M1$ mixing ratio, the F_ν coefficients are defined by

$$F_\nu(LL'I'I) = (-1)^{I'+3I-1}[(2I+1)(2L+1)(2L'+1)]^{\frac{1}{2}}C(LL'\nu;1-10)W(LL'II;\nu I'), \tag{4}$$

C is the Clebsch-Gordan coefficient $C(J_1J_2J_3;M_1M_2M_3)$ and W is the Racah coefficient. The parity conserving (PC), quantity is given by [23,24]

$$R_\gamma^{\rm PC}(\cos\theta) = \left(\sum_{\nu=1,3} P_\nu(\cos\theta)B_\nu(2)[F_\nu(1122) + F_\nu(2222)\delta_-^2 + 2F_\nu(1222)\delta_-]\right)$$
$$\times\left(\sum_{\nu=0,2,4} P_\nu(\cos\theta)B_\nu(2)[F_\nu(1122) + F_\nu(2222)\delta_-^2 + 2F_\nu(1222)\delta_-]\right)^{-1}, \tag{5}$$

where

$$B_\nu(2) = \sum_M (2\nu+1)^{\frac{1}{2}}C(2\nu 2;M0M)p(M). \tag{6}$$

$p(M)$ is the polarization fraction of the M state, which determines the degree of the orientation of the nucleus.

In order to measure a PNC effect one must find situations for which the $R_\gamma^{\rm PC}$ part in Eq. (1) vanishes. Two particular cases have this property: (i) The case of an initially unpolarized nucleus for which $B_0(2) = 1$, $B_{\nu\neq 0}(2) = 0$, and $F_0(LL'22) = \delta_{LL'}$. In this particularly simple case P_γ reduces to the well-known expression of the circular polarization, $(P_\gamma)_0$. (ii) One may prepare a polarized state by choosing $p(M) = \delta_{M0}$ for which, $B_{\nu=1,3}(2) = 0$ and the $R_\gamma^{\rm PC}$ part vanishes.

Another observable which measures a PNC effect is the forward-backward asymmetry of the gamma rays emitted by polarized nuclei

$$A_\gamma(\theta) \equiv \frac{W(\theta) - W(\pi-\theta)}{W(\theta) + W(\pi-\theta)}. \tag{7}$$

This observable has been successfully used in the ^{19}F case [9,10] in order to avoid the small efficiency of the

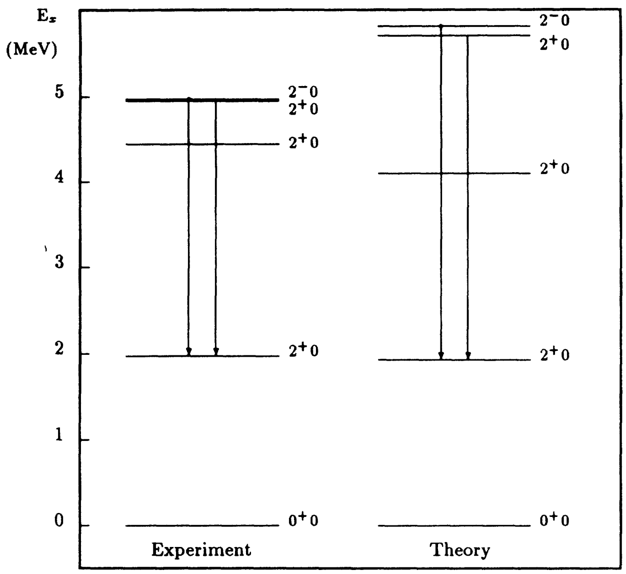

FIG. 2. Same as Fig. 1 for the low $2^{\pm}0$ levels in ^{36}Ar.

Compton polarimeters when one measures the degree of circular polarization. If the mixing ratios are small ($\delta_+, \delta_- \ll 1$) one can show that [21]

$$A_\gamma(\theta) \simeq (P_\gamma)_0 R_\gamma^{\rm PC}(\cos\theta) \ , \qquad (8)$$

where θ represents the angle between the emitted photon and the axis of polarization (if any). The angular distribution described by this formula has a maximum for $\theta = 0°$ [21]. It has the advantage that the parity conserving (PC) circular polarization, $R_\gamma^{\rm PC}(\theta)$ in Eq. (8), can be measured experimentally. $(P_\gamma)_0$ is the essential quantity for all PNC observables.

In order to determine the magnitude of $(P_\gamma)_0$ we have made a shell-model estimate of the PNC matrix element

$$M_{\rm PNC} = \langle J^\pi T, E_x({\rm MeV}) \mid H_{\rm PNC} \mid J^{-\pi} T' E_x'({\rm MeV})\rangle, \qquad (9)$$

where $H_{\rm PNC}$ is the PNC Hamiltonian given by Desplanques, Donoghue, and Holstein (DDH) [12], Dubovik and Zenkin (DZ) [18], Adelberger and Haxton (AH) [11], or Kaiser and Meissner (KM) [19].

The calculations were carried out with the shell-model code OXBASH [25] in the $2s1d$-$2p1f$ model space in which the $2s_{1/2}, 1d_{5/2}, 1d_{3/2}, 2p_{1/2}, 2p_{3/2}, 1f_{7/2}$, and $1f_{5/2}$ orbitals are active. The truncations we made within this model space were $(2s1d)^{20}$ $(0\hbar\omega)$ for the positive parity states and $(2s1d)^{19}(2p1f)^1$ $(1\hbar\omega)$ for the negative parity states. These truncations are necessary due to the dimension limitations, but we believe that they are realistic. The Brown-Wildenthal interaction [26] was used for the positive parity states and the Warburton-Becker-Brown interaction [22] was used for the negative parity states. Both interactions have been tested extensively with regards to their reproduction of spectroscopic properties [22,26]. The calculation of the PNC matrix element which included both the core (inactive) and active orbitals has been performed as described in Ref. [27].

All the components [12,11] of the parity nonconserving potential are short-range two-body operators. Because the behavior of the shell-model wave functions at small NN distances has to be modified, short-range correlations (SRC) were included by multiplying the harmonic oscillator wave functions (with $\hbar\omega = 45A^{-\frac{1}{3}}$ MeV $- 25A^{-\frac{2}{3}}$ MeV) by the Miller and Spencer factor [28]. This procedure is consistent with results obtained by using more elaborate treatments of SRC such as the generalized Bethe-Goldstone approach [7,8]. The PNC pion exchange matrix is decreased by 30–50% as compared with the values of the matrix elements without including SRC, while the ρ (ω) exchange matrix elements is much smaller (by a factor of $\frac{1}{3}$ to $\frac{1}{6}$).

The calculated excitation energies of the first three $T = 0$, 2^+ levels in ^{36}Ar are 1.927, 4.410, and 7.174 MeV. The first two of these are in good agreement with experimental levels at 1.970 and 4.440 MeV. The third experimental 2^+ state at E_x=4.951 MeV (the state belonging to the parity doublet) apparently is an intruder state in the $2s1d$ $(0\hbar\omega)$ model space. This conclusion is also supported by the suppressed Gamow-Teller β transition probability to this third state [29]. We have thus expanded the model space to include some $2\hbar\omega$ configurations — those of the type $(1d_{5/2})^{12}(2s_{1/2}, 1d_{3/2})^6(2p_{3/2}, 1f_{7/2}, 2p_{1/2})^2$. The $2\hbar\omega$ configurations were shifted down by 11.5 MeV relative to the $0\hbar\omega$ configurations, so that the first 2^+0 state with a dominant $2\hbar\omega$ component ($\sim 80\%$) becomes the third 2^+ in the calculated $(0+2)\hbar\omega$ spectrum. The dominant PNC transition is $1d_{3/2}$ - $2p_{3/2}$ and the DDH PNC matrix elements is 0.12 eV (see Table I). Due to the truncations made, the PNC result for ^{36}Ar may not be as reliable as that for ^{36}Cl.

The calculated excitation energies of the first three $T = 1$, 2^+ levels in ^{36}Cl are 0, 2.008, 2.451, and 4.429 MeV. The first three of these are in good agreement with experimental levels at 0, 1.959, and 2.492 MeV. The theoretical $B(E2)$ and $B(M1)$ and mixing ratios are in relatively good agreement with the experiment (see Table I). For both ^{36}Ar and ^{36}Cl the 2^- states are the lowest observed experimentally and the theoretical wave functions should thus be fairly reliable. The extremely weak $E1$ transitions do not serve as a useful test of the wave functions (the one in ^{36}Ar is isospin forbidden).

DDH [12] investigated a variety of approximations within the quark model for the weak coupling constants and discussed the model uncertainties. These uncertainties give rise to a range of values for the PNC coupling constants. Recently several other calculations have been made, one within the framework of the chiral soliton model [19] and others within the quark framework [18,20]. Even though both of these approaches lead to fixed values for PNC coupling constants [19] (see Table II), the values are subject to uncertainties. In particular, the soliton model gives an extremely small value for h_π as compared to DDH, however, this result comes essentially

TABLE I. Input data, physical quantities, and theoretical PNC matrix elements necessary for calculating γ circular polarizations and asymmetries for the two PMD cases studied in the present work. The experimental data is taken from Ref. [30] unless noted.

	^{36}Cl	^{36}Ar
$I_i^\pi T_i, E_i$ (MeV) $\rightarrow$ $I_f^\pi T_f, E_f$ (MeV)	$2^+1, 1.959$ MeV $\rightarrow$ 2^+1, g.s.	$2^+0, 4.951$ MeV $\rightarrow$ $2^+0, 1.97$ MeV
Lifetime (τ_+)	60 $\pm$ 15 fs	$\leq$ 50 fs
Branching ratio (b_+)	94.4 %	15 %
Mixing ratio $(\delta_+)_{\rm expt}$	(-5.2 $\pm$ 1.6) or (- 0.19 $\pm$ 0.06) [31]	
Mixing ratio $(\delta_+)_{\rm theor}$	-0.24	0.41
$B(M1)_{\rm expt}(\mu_N^2)$	0.08 (δ_+=-0.2);0.003 (δ_+=-5.2)	
$B(M1)_{\rm theor}(\mu_N^2)$	0.14	0.0009
$B(E2)_{\rm expt}(e^2$ fm$^4)$	12 (δ_+=-0.2); 298 (δ_+=-5.2)	
$B(E2)_{\rm theor}(e^2$ fm$^4)$	30	0.27
$I_i^\pi T_i, E_i$ (MeV) $\rightarrow$ $I_f^\pi T_f, E_f$ (MeV)	$2^-1, 1.951$ MeV $\rightarrow$ 2^+1, g.s.	$2^-0, 4.974$ MeV $\rightarrow$ $2^+0, 1.97$ MeV
Lifetime (τ_-)	2.6 $\pm$ 0.3 ps	14 $\pm$ 5 ps
Branching ratio (b_-)	60 %	4.0 $\pm$ 0.9 %
Mixing ratio $(\delta_-)_{\rm expt}$	(- 0.10 $\pm$ 0.10) [31]	
Mixing ratio $(\delta_-)_{\rm theor}$	0.009	
$B(E1)_{\rm expt}(e^2$ fm$^2)$	$1.9\cdot10^{-5}$	0.7×10^{-7} (if $\delta_-=0.0$)
$B(E1)_{\rm theor}(e^2$ fm$^2)$	0.008	
$B(M2)_{\rm expt}(\mu_N^2$ fm$^2)$	$\leq$ 25	
$B(M2)_{\rm theor}(\mu_N^2$ fm$^2)$	2.5	0.24
$M_{\rm PNC}^{\rm DDH}$ (eV)	-0.019	0.122
$M_{\rm PNC}^{\rm DDH}$ (eV), $h_\pi^1=\frac{1}{4}(h_\pi^1)_{\rm DDH}$	-0.057	0.122
$M_{\rm PNC}^{\rm KM}$ (eV)	-0.023	0.067

from the factorization approximation, and DDH discuss the importance of going beyond the factorization approximation [12]. In any case it is clear that the observation of PNC in nuclei is a test not only of the $\Delta-S=0$ PNC component of the weak interaction, but also of the hadronic strong interaction models.

The results (up to a complex phase factor) can be summarized as

$$M_{\rm PNC}(^{36}{\rm Cl}) = (1.09h_\pi^1 - 0.20h_\rho^1 - 0.30h_\omega^1 - 0.027h_{\rho'}^1 + 0.57h_\rho^0 + 0.32h_\omega^0 + 0.015h_\rho^2)\times10^{-2}{\rm eV} \qquad (10)$$

and

$$M_{\rm PNC}(^{36}{\rm Ar}) = -(1.00h_\rho^0 + 0.44h_\omega^0)\times10^{-2}\ {\rm eV}\ . \qquad (11)$$

Here $h_{\rm meson}^{\Delta T}$ should be given in units of 10^{-7} as in Table II.

In the ^{36}Cl case the components containing h_π (M_π) and $h_{\rho(\omega)}$ $(M_{\rho+\omega})$ couplings come in with opposite signs, however the difference is remarkably almost the same in all the weak interaction models. The specific numbers are $M_\pi=0.050$ eV and $M_{\rho+\omega}=-0.069$ eV for DDH, and $M_\pi=0.002\,07$ eV and $M_{\rho+\omega}=-0.023$ eV for KM. In the ^{36}Ar case the h_ρ components strongly dominate the PNC matrix element ($M_{\rm PNC}$) (e.g., within DDH: $M_\rho \doteq 0.113$ eV, while $M_\omega=0.009$ eV).

The PNC matrix elements obtained are a factor of 3–6 smaller than the typical "isoscalar" matrix elements in $A=14$–20 nuclei (e.g., ~0.3 eV in ^{19}F [11] and ^{14}N [32]). An analysis of the different contributions to the PNC

TABLE II. Weak meson-nucleon coupling constants calculated within different weak interaction models (in units of 10^{-7}). The abbreviations are KM, Kaiser and Meissner [19]; DDH, Desplanques, Donoghue, and Holstein [12]; AH, Adelberger and Haxton [11]; and DZ, Dubovik and Zenkin [18]. The $g_{\rm meson}$'s needed to obtain these results were taken from Ref. [11].

$h_{\rm meson}^{\Delta T}$	KM	DDH	AH	DZ
h_π^1	0.19	4.54	2.09	1.30
h_ρ^0	-3.70	-11.40	-5.77	-8.30
h_ρ^1	-0.10	-0.19	-0.22	0.39
h_ρ^2	-3.30	-9.50	-7.06	-6.70
$h_{\rho'}^1$	-2.20	0.00	0.00	0.00
h_ω^0	-1.40	-1.90	-4.97	-3.90
h_ω^1	-1.00	-1.10	-2.39	-2.20

matrix elements indicates the reasons for this behavior.

(i) The one-body transition densities (OBTD) for the isoscalar and isovector $1d_{3/2}$-$2p_{3/2}$, $2s_{1/2}$-$2p_{1/2}$, $1d_{5/2}$-$1f_{5/2}$ PNC transitions are a factor 5–10 smaller than the dominant $1p_{1/2}$-$2s_{1/2}$ transition in the $A = 14$–20 nuclei. This can be understood by the fact that for $A = 36$ nuclei the most probable subshell to be filled in the $1f2p$ major shell is $1f_{7/2}$ which is 2 MeV lower than $2p_{3/2}$. This substantially decreases the occupation probability of the $2p_{3/2}, 2p_{1/2}$, and $1f_{5/2}$ states which are important for PNC transitions. In the $A = 14$–20 region the $2s_{1/2}$ is nearly degenerate with the $1d_{5/2}$ state and has a higher occupation probability. In the ^{36}Cl case the dominant OBTD is for the transition $1d_{3/2}$-$2p_{3/2}$. We have estimated the $2\hbar\omega$–$1\hbar\omega$ and $2\hbar\omega$–$3\hbar\omega$ contribution to this dominant transition within the $(2s_{1/2}, 1d_{3/2}, 2p_{3/2}, 1f_{7/2})^8$ mode space. The $2\hbar\omega$–$1\hbar\omega$ and $2\hbar\omega$–$3\hbar\omega$ contribution have a magnitude of 5–7 % of the $0\hbar\omega$–$1\hbar\omega$ and are opposite in sign, so they are small and also cancel each other. These calculations suggest that the $0\hbar\omega$–$1\hbar\omega$ contribution is the dominant term. This behavior is in contrast with that in $A = 18$–21 nuclei, for which the higher $n\hbar\omega$ contributions seems to be important [33]. One can understand this by the relatively weaker coupling between the sd and fp shells as compared with the coupling between the p and sd shells.

(ii) In the case of ^{36}Cl, the isoscalar and isovector contributions have opposite signs leading to a large suppression of the total PNC matrix element (for the DDH weak coupling constants). Considering the strong constraints on h^1_π given by the ^{18}F experiments, i.e. $h^1_\pi \leq (1/4)(h^1_\pi)_{\rm DDH}$ [34], one obtains for $(M^{^{36}Cl}_{\rm PNC})_{\rm DDH'}$ a value of -0.057 eV.

Taking into account the results of the above discussion, we have used $M^{^{36}\rm Cl}_{\rm PNC} = -0.057$ eV and $M^{^{36}\rm Ar}_{\rm PNC} = 0.122$ eV to calculate the $(P_\gamma)_0$. We obtain $(P_\gamma)_0^{^{36}\rm Cl} = 1.1 \times 10^{-4}$ and $(P_\gamma)_0^{^{36}\rm Ar} = 3.5 \times 10^{-4}$ (we used the small mixing ratios, which agree with our calculations). These values can be favorably compared with the experimental upper limit, $P_\gamma \leq 3.9 \times 10^{-4}$, for ^{18}F [34].

One can try to avoid the low efficiency of the Compton polarimeters by measuring the forward-backward asymmetry, Eqs. (7) and (8). In this case one must find an efficient polarization transfer mechanism which would permit one to obtain a PC circular polarization $R^{\rm PC}_\gamma$ larger than the polarimeter efficiency ($\sim 1\%$ [11]). For example, the ^{36}Ar PMD can be populated in the ^{39}K $(\vec{p}, \alpha)^{36}$Ar reaction (in analogy with the ^{19}F case [9,10]), with low energy ($E_p \geq 3.7$ MeV) protons, while the ^{36}Cl PMD can be populated in the ^{39}K $(\vec{n}, \alpha)^{36}$Cl reaction with relatively slow ($E_n \geq 0.6$ MeV) neutrons.

In conclusion, we have theoretically analyzed two new PMD cases in mass $A = 36$ nuclei. The parity nonconserving transition for the PMD in ^{36}Ar is isoscalar and the corresponding PNC observable is sensitive to the dominant h^0_ρ weak coupling constant, analogous to the 0^+1, 0^-1 doublet in ^{14}N [32]. The parity nonconserving transition for the PMD in ^{36}Cl is isoscalar+isovector and the corresponding PNC observable is sensitive to the combination of h^0_ρ and h^1_π. From this point of view it is analogous to the ^{19}F case. However, the possible information extracted from this case could be complementary to the ^{19}F result. Here, the isoscalar and isovector contributions are out of phase, while for ^{19}F they are in phase. The predicted circular polarizations of the order of magnitude 10^{-4} are within the limits of accuracy of the existent experimental setups.

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