P and T odd nuclear moments

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Abstract

We discuss the nuclear anapole moment, a new P odd, T even electromagnetic multipole. Its discovery in the cesium experiment gives a first-rate information on parity-nonconserving nuclear forces. New prospects in the field are considered. Upper limits on the electric dipole moments, which are P odd, T odd multipoles, of elementary particles and atoms are presented, and their physical implications discussed. The atomic experiments are demonstrated to be as informative in this respect as the neutron ones. Tremendous progress in the field can be expected from experiments at ion storage rings and linear electrostatic traps.

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1 Introduction

We discuss in the present article two sorts of electromagnetic nuclear moments. To the first type belong those which violate the invariance under P, the parity transformation, but conserve T, time reversal invariance. To the second type belong the moments violating both P and T.

P odd, T even moments arise in a system without centre of inversion. Though their existence was predicted in the middle of the past century [1], the lowest of them, the so-called anapole, that of the $^{133}$Cs nucleus, was experimentally discovered only few years ago [2]. Being of a great interest by itself, this discovery gives also quite interesting and nontrivial information on P odd nuclear forces.

As to P odd, T odd multipoles, the lowest of them, electric dipole moments (EDM), have interested physicists since 1950, when it was first suggested that there was no experimental evidence that nuclear forces are symmetric under parity transformation [3].

In 1964 it was discovered that the invariance under CP transformation, which combines charge conjugation with parity, is violated in $K$-meson decays. This provided a new incentive for EDM searches. Since the combined operations of CPT are expected to leave a system invariant, breakdown of CP invariance should be accompanied by T violation. Thus there is a reason to expect that EDMs should exist at some level.

The original neutron EDM experiments were later supplemented with searches for EDMs of other objects, nuclei included. These investigations are pursued now by many groups. Over the years, the upper limit on the neutron EDM has been improved by seven orders of magnitude, and the upper limit on the electron EDM obtained in atomic experiments is even more strict.

Even without the discovery of the effect sought, the neutron and atomic experiments have ruled out most models of CP violation. As to the mechanism of CP violation incorporated in the standard model of electroweak interactions, which is most popular at present, the predictions for the neutron and nuclear EDMs are roughly six orders of magnitude below the present experimental bound. The gap for the electron EDM is much larger.

But does this mean that the EDM experiments are of no serious interest for the elementary particle physics, that they are nothing but mere exercises in precision spectroscopy? Just the opposite. It means that the EDM experiments now, at the present level of accuracy, are extremely sensitive to possible new physics beyond the standard model, physics to which the kaon decays are insensitive.

2 Nuclear anapole moments

2.1 General discussion. Anapole moment of $^{133}$Cs

In a system which has no definite parity, a special distribution of magnetic field may arise [1]. It cannot be reduced to common electromagnetic multipoles, such as dipole or quadrupole moments, but looks like the magnetic field created by a current in toroidal
winding. This special source of electromagnetic field was called (by A.S. Kompaneets) anapole.

For many years the anapole remained a theoretical curiosity only. The situation has changed with the investigations of parity nonconservation (PNC) in atoms. The tiny P odd effects are enhanced therein in particular by the weak nuclear charge $Q_W$, which is close numerically to the neutron number and is thus about a hundred in heavy atoms. However, this enhancement refers only to the nuclear-spin-independent weak interaction. Meanwhile, the nuclear-spin-dependent effects due to neutral currents not only lack the mentioned coherent enhancement, but are also strongly suppressed numerically in the electroweak theory. Therefore, the observation of PNC nuclear-spin-dependent effects in atoms looked absolutely unrealistic.

However, it was demonstrated [4, 5] that these effects in heavy atoms are dominated not by the weak interaction of neutral currents, but by the electromagnetic interaction of atomic electrons with nuclear anapole moment (AM). It should be mentioned first of all that the magnetic field of an anapole is contained within it, in the same way as the magnetic field of a toroidal winding is completely confined inside the winding. It means that the electromagnetic interaction of an electron with the nuclear AM occurs only as long as the electron wave function penetrates the nucleus. In other words, this electromagnetic interaction is as local as the weak interaction, they are quite similar. The nuclear AM is induced by PNC nuclear forces and is therefore proportional to the same Fermi constant $G = 1.027 \times 10^{-5}m^{-2}$ (we use the units $\hbar = 1, c = 1; m$ is the proton mass), which determines the magnitude of the weak interactions in general and that of neutral currents in particular. The electron interaction with the AM, being of the electromagnetic nature, introduces an extra small factor into the effect discussed, the fine-structure constant $\alpha = 1/137$. Then, how it comes that this effect is dominating? The answer follows from the same picture of a toroidal winding. It is only natural that the interaction discussed is proportional to the magnetic flux through such a winding, and hence in our case to the cross-section of the nucleus, i.e. to $A^{2/3}$, where $A$ is the atomic number. In heavy nuclei this enhancement factor is close to 30 and compensates essentially for the smallness of the fine-structure constant $\alpha$. As a result, the dimensionless effective constant $\kappa$ which characterizes the anapole interaction in the units of $G$ is not so small in heavy atoms, but is numerically close to 0.3 (we use the same definition of the effective constant $\kappa$ as in [4, 5]).

Still, the interaction discussed constitutes only about one percent of the main atomic PNC effect, which is due to $Q_W$. To single out the anapole interaction one should compare the PNC effects for different hyperfine components of an optical transition. The main effect, independent of the nuclear spin, will obviously be the same for all components. But the anapole contribution depends on the mutual orientation of the nuclear spin and the electron total angular momentum, and thus changes from one hyperfine component to another. The observation of this tiny effect is an extremely difficult problem, and it is no accident that the searches for the nuclear AM demanded many years of hard work by several groups.

The nuclear anapole moment was experimentally discovered in 1997 [2]. This result for
the total effective constant of the PNC nuclear-spin-dependent interaction in $^{137}$Cs is

$$\kappa_{\text{tot}} = 0.44(6).$$  \hspace{1cm} (1)

To extract this number from experimental data, the results of atomic calculations \[6, 7\] were used; these calculations were performed in different approaches, but are in excellent agreement. Their accuracy is no worse than 2 – 3%. If one excludes the neutral current nuclear-spin-dependent contribution from the above number, as well as the result of the combined action of $Q_W$ and the usual hyperfine interaction, the answer for the anapole constant will be

$$\kappa = 0.37(6).$$  \hspace{1cm} (2)

Thus, the existence of an AM of the $^{137}$Cs nucleus is reliably established. A beautiful new physical phenomenon, a peculiar electromagnetic multipole has been discovered.

### 2.2 Theoretical issues. Implications for P odd nuclear forces

But the discussed result does not reduce to only this. It brings valuable information on PNC nuclear forces. Of course, to this end it should be combined with reliable nuclear calculations. However, it is instructive to start, as it was done in \[5\], with a rather crude approximation. Not only one assumes here that the nuclear spin $I$ coincides with the total angular momentum of an odd valence nucleon, while the other nucleons form a core with the zero angular momentum. The next assumption is that the core density $\rho(r)$ is constant throughout the space and coincides with the mean nuclear density $\rho_0$. The last assumption is reasonable if the wave function of the external nucleon is mainly localized in the region of the core. Then simple calculations give the following result for the anapole constant \[5\]:

$$\kappa = \frac{9}{10} g \frac{\alpha \mu}{m r_0^2} A^{2/3}.\hspace{1cm} (3)$$

Here $g$ is the effective constant of the P odd interaction of the outer nucleon with the nuclear core, $\mu$ is the magnetic moment of the outer nucleon, $r_0 = 1.2$ fm. The so-called “best values” for the parameters of P odd nuclear forces \[8\] result in $g_p = 4.5$ for an outer proton \[5, 9, 10\]. The $A$-dependence of this constant has been already anticipated from the picture of a toroidal winding.

Various theoretical predictions (with $g_p = 4.5$) for the AMs of nuclei of those atoms where the PNC effects have been studied experimentally, are presented in Table 1 (below in this subsection we follow \[11\]). The analytical estimate \[3\] produces smooth $A^{2/3}$ behaviour (see the first column of Table 1), but certainly exaggerates the effect due to the assumption that the P odd contact interaction with the nuclear core extends throughout the whole localization region of the unpaired nucleon. Indeed, more refined calculations, already in the single-particle approximation (SPA) (in the second column of Table 1 we present the results obtained with the Woods-Saxon potential, including contributions of contact and spin-orbit currents, which is perhaps the most advanced calculation in SPA) reveal certain shell effects quite pronounced in the values of $\kappa$ for Tl and Bi. Both these nuclei are close...
to the doubly-magic $^{208}$Pb. However, while the anapole moment of Tl nucleus in SPA is close to its analytical estimate, the anapole moment of Bi in SPA differs significantly both from the analytical formula and from the anapole moment of Tl. This difference can be attributed to the difference in the single-particle orbitals for the unpaired proton in Tl and Bi. The $3s_{1/2}$ wave function in Tl is concentrated essentially inside the nuclear core, while the $1h_{9/2}$ wave function in Bi is pushed strongly outside of it. By this reason the unpaired proton in Bi “feels” in fact much smaller part of the P odd weak potential. An analogous suppression of the PNC interaction takes place for the outer $1g_{7/2}$ proton in Cs.

Various approaches were used as well in the many-body calculations. In one of them [14] the random-phase approximation (RPA) was employed to calculate the effects of the core polarization. In another approach [15] large basis shell-model calculations were performed. However, in the last case there is a serious problem: the basis necessary to describe simultaneously the effects of both regular nuclear forces and P odd ones, is in fact too large. Therefore, some additional approximations were made in [15] in order to reduce the size of the basis space.

Fortunately, the Tl nucleus is a rather special case in the many-body approach as well. Not only is it close to the doubly-magic $^{208}$Pb, but its unpaired proton is $3s_{1/2}$, but not $1h_{9/2}$ as in Bi. This makes the effects of the core polarization here relatively small. Thus the density of states in Tl is reduced, and an effective Hamiltonian suitable for shell-model calculations can be constructed [16]. This Hamiltonian was used in [13] to calculate the anapole moment of Tl nucleus. The result of [13] and the RPA result of [14] for the thallium coincide, in spite of completely different descriptions of nuclear forces used in these works to calculate the core polarization. These results of [13, 14] differ essentially from the value obtained in [15] under extra assumptions: the closure approximation and further reduction of a three-body matrix element to the two-body one. It is also worth mentioning perhaps that in [13, 14] and [15] different parameterizations of the parity violating nuclear forces have been used.

|     | 5   | 12  | 13  | 14  | 15  |
|-----|-----|-----|-----|-----|-----|
| $^{133}$Cs | 0.37 | 0.22 | —   | 0.15 | 0.21 |
| $^{203,205}$Tl | 0.49 | 0.37 | 0.24 | 0.24 | 0.10 |
| $^{209}$Bi   | 0.51 | 0.30 | —   | 0.15 | —   |

Table 1
Thus we believe that the theoretical predictions for the AMs of nuclei of the present experimental interest, can be reasonably summarized now, at “best values” of P odd constants, as follows:

$$\kappa^{(133\text{Cs})} = 0.15 - 0.21, \quad \kappa^{(203,205\text{Tl})} = 0.24, \quad \kappa^{(209\text{Bi})} = 0.15.$$  \hspace{1cm} (4)

We believe also that there are good reasons to consider these predictions as sufficiently reliable, at the accepted values of the P odd nuclear constants.

The comparison of the value (4) for the cesium AM with the experimental result (2) indicates that the “best values” of [8] somewhat underestimate the magnitude of P odd nuclear forces. In no way is this conclusion trivial. The point is that the magnitude of parity-nonconserving effects found in some nuclear experiments is much smaller than that following from the “best values” (see review [17]). In all these experiments, however, either the experimental accuracy is not high enough, or the theoretical interpretation is not sufficiently convincing. The experiment [2] looks much more reliable in both respects. Therefore, in line with its general physics interest, the investigation of nuclear AMs in atomic experiments is first-rate, almost table-top nuclear physics.

It is appropriate perhaps to point out here a problem we still have not mentioned. The experimental result for the thallium AM, $$\kappa = -0.22 \pm 0.30$$ [18], does not comply with the theoretical prediction for it presented in (4) (the disagreement will be even more serious if one assumes that the nuclear P odd constants are larger than the “best values” of [8] as indicated by the measurement of the cesium AM). Obviously, it is highly desirable for this problem to be cleared up.

2.3 Prospects

The experiments aimed at the detection and measurement of nuclear AMs are extremely difficult. Therefore, it would be quite important to find a situation where the AM effects are strongly enhanced. Unfortunately, up to now nobody could point out any example of a pronounced enhancement of a nuclear AM by itself.

In principle, the AM can be enhanced in the case when anomalously close to the ground state of a nucleus there is an opposite-parity level of the same angular momentum. In this connection, attention was attracted in [19, 20] to exotic halo nuclei. In particular, the exotic neutron-rich halo nucleus $^{11}\text{Be}$ was considered therein. In this nucleus the outer odd neutron is in the state $2s_{1/2}$, its only bound excited level being $1p_{1/2}$ (the well-known “inversion of levels”). The anomalously small energy separation between these two levels of opposite parity,

$$|\Delta E| = E(1p_{1/2}) - E(2s_{1/2}) = 0.32 \text{ MeV},$$  \hspace{1cm} (5)

enhances by itself their P-odd mixing and thus the AM of this nucleus. As pointed out in [19, 20], the small binding energy of the odd neutron,

$$|\Delta E_0| = 0.50 \text{ MeV},$$  \hspace{1cm} (6)
affects the AM additionally, but in two opposite directions. On one hand, it suppresses the overlap of the odd-neutron wave function with the core, and thus suppresses the mixing of the \(2s_{1/2}\) and \(1p_{1/2}\) levels due to the weak interaction operator which looks as

\[
W = \frac{G}{\sqrt{2}} \frac{g_n}{2m} \{ \sigma \mathbf{p}, \rho(r) \};
\]

(7)

here \(g_n\) is the effective constant of the P-odd interaction of the outer neutron with the nuclear core, \(\sigma\) and \(\mathbf{p}\) are the spin and momentum operators of the outer neutron, and \(\rho(r)\) is the spherically symmetric core density. On the other hand, the small binding energy enhances the matrix element of \(\mathbf{r}\) in the anapole operator of the neutron

\[
\mathbf{a} = \frac{\pi e \mu_n}{m} \mathbf{r} \times \sigma;
\]

(8)

here \(\mu_n = -1.91\) is the neutron magnetic moment.

The detailed calculation which takes into account the P-odd mixing of the ground state with the \(1p_{1/2}\) level only, results in the following value for the effective anapole constant \([20]\):

\[
\kappa_1(\text{^{11}Be}) = 0.17 g_n.
\]

(9)

Indeed, this value is 15 times larger than that given by the estimate \([3]\) for \(A = 11\) (the neutron constant \(g_n\) is poorly known by itself, most probably \(g_n \lesssim 1\)). Certainly, this enhancement of an AM in a light nucleus would be of a serious interest, even if its possible experimental implications are set aside.

However, a strong enhancement of AM, as given in \([3]\), in a loosely bound nucleus does not look natural. In particular, nothing of the kind happens in the deuteron. Even in the limit of vanishing binding energy, when the energy interval between the deuteron \(s\) state and the continuum \(p\) states tends to zero, the deuteron AM is not enhanced essentially (see below). As to the problem of \(^{11}\)Be discussed here, a strong cancellation between the contribution of the bound \(1p_{1/2}\) state (accounted for in \([3]\)) and that of the continuum (omitted therein) takes place \([21]\). This cancellation results in a serious suppression of the naive estimate as compared to \([3]\):

\[
\kappa(\text{^{11}Be}) \simeq 0.07 g_n.
\]

(10)

Let us come back now to the above mentioned case of the deuteron AM. It can be calculated in a closed form in the chiral limit, i. e. for the vanishing pion mass, \(m_\pi \rightarrow 0\) \([22]\) (see also \([23]\)). In this limit one can confine for the loosely bound deuteron to the weak one-pion exchange. We use the Lagrangians of the strong \(\pi NN\) interaction and of the weak P-odd one, \(L_s\) and \(L_w\), in the form

\[
L_s = g \left[ \sqrt{2} (\bar{p}i\gamma_5 n \pi^+ + \bar{n}i\gamma_5 p \pi^-) + (\bar{p}i\gamma_5 p - \bar{n}i\gamma_5 n) \pi^0 \right];
\]

(11)

\[
L_w = \bar{g} \sqrt{2} i (\bar{p} n \pi^+ - \bar{n} p \pi^-).
\]

(12)
Then in the zero-range approximation for the deuteron wave function, one obtains the following result for its AM

\[ a_d^{(0)} = - \frac{eg\bar{g}}{6mm_\pi} \frac{1 + \xi}{(1 + 2\xi)^2} \left( \mu_p - \mu_n - \frac{1}{3} \right) I. \]  

(13)

Here \( \mu_p = 2.79 \) and \( \mu_n = -1.91 \) are the proton and neutron magnetic moments, respectively; \( I \) is the deuteron spin. With the deuteron binding energy \( \varepsilon = 2.23 \) MeV, the numerical value of the parameter \( \xi = \sqrt{\varepsilon m/m_\pi} \) in this formula is \( \xi = 0.32 \). As it was mentioned above, even for the vanishing binding energy \( \varepsilon \to 0 \) no essential enhancement occurs for the deuteron AM.

Let us note that in the same chiral limit the nucleon AM can be calculated in a closed form. It was done in 1980 by I.B. Khriplovich and A.I. Vainshtein. The result is the same for the proton and neutron:

\[ a_p = a_n = - \frac{eg\bar{g}}{12mm_\pi} \left( 1 - \frac{6}{\pi} \frac{m_\pi}{m} \ln \frac{m}{m_\pi} \right) \sigma. \]  

(14)

To obtain the total deuteron AM, one should combine the contributions of the nucleon anapoles (14) with (13).

To summarize, no example of a considerable enhancement of nuclear AMs by internal nuclear effects has been found up to now.

There is however a situation when the manifestation of a nuclear AM is strongly enhanced by atomic effects. We mean a proposal (which looks at the moment the most promising one) to measure PNC in the strongly forbidden \( 6s^2 \frac{1}{2}S_0 \to 6s5d \frac{3}{2}D_1 \) transition.
in ytterbium [24]. The advantage of this transition is that the PNC effect in it is more than 100 times larger than in cesium. The enhancement is due to the fact that the $^3D_1$ state is close ($\Delta E \simeq 600 \text{ cm}^{-1}$) to a level of opposite parity $^1P_1$ (see Figure 1) whose composition is $6s6p$ with a strong admixture of $5d6p$. Due to this $5d6p$ component, there is large PNC mixing between $^3D_1$ and $^1P_1$. The relatively simple atomic spectrum of Yb allows one to perform atomic calculations of the PNC effect with an accuracy about 20% [24–26].

Ytterbium has seven stable isotopes between A=168 and A=176. Two of them, $^{171}$Yb, $I=1/2$ and $^{173}$Yb, $I=5/2$, with non-zero nuclear spin, can be used to measure the AMs. With a valence neutron in these nuclei, such measurements will be a valuable complement to the cesium anapole result in the determination of the PNC nuclear constants.

Moreover, one more transition in Yb, $6s^2 \, ^1S_0 \rightarrow 6s5d \, ^3D_2$, is of a special interest for the anapole measurements. PNC mixing between $^3D_2$ and $^1P_1$ states (their separation is $\Delta E \simeq 300 \text{ cm}^{-1}$) is only due to that P odd interaction of electrons with nuclear spin which possess $\Delta J = 1$. Thus, in this transition the anapole interaction will be the main source of P odd effects, rather than a small correction to the dominant nuclear-spin-independent interaction as it is the case with $^1S_0 \rightarrow ^3D_1$.

Some preliminary spectroscopic measurements in ytterbium, related to the discussed experiments, have been done. They resulted in the lifetime of the $^3D_1$ state, $380(30)$ ns [27], as well as in the values of the E2 amplitude of the $^1S_0 \rightarrow ^3D_2$ transition, $0.65(3) \text{ ea}^2$ [28], and of the strongly forbidden M1 amplitude of the $^1S_0 \rightarrow ^3D_1$ transition, $1.33(26) \times 10^{-4} \mu_B$ [29].

3 Nuclear electric dipole moments

We start here with the present experimental information on the EDMs.

3.1 Elementary particles

The experimental upper limit on the neutron EDM is [30, 31, 32]

$$d_n < (6 - 10) \times 10^{-26} \text{ e cm}.$$  \hspace{1cm} (15)

The sensitivity of these experiments can be, hopefully, improved by 2 – 3 orders of magnitude.

The best result for the electron EDM

$$d_e = (0.69 \pm 0.74) \times 10^{-27} \text{ e cm}$$  \hspace{1cm} (16)

was obtained in atomic experiment with Tl [33]. Hopefully, this limit can be pushed well into the $10^{-28}$ e cm range.

I would like to quote here one more upper limit, that on the muon EDM [34]:

$$d_\mu < 10^{-18} \text{ e cm}.$$  \hspace{1cm} (17)

An experiment was recently proposed to search for the muon EDM with the sensitivity of $10^{-24}$ e cm [35]. We will come back to this proposal later.
The predictions of the Standard Model are, respectively:

\[ d_n \sim 10^{-32} - 10^{-31} \text{ e cm}; \]  
\[ d_e < 10^{-40} \text{ e cm}; \]  
\[ d_\mu < 10^{-38} \text{ e cm}. \] (18) (19) (20)

### 3.2 Atoms and Nuclei

The best upper limit on EDM of anything was obtained in atomic experiment with $^{199}$Hg [36]. The result for the dipole moment of this atom is

\[ d(^{199}\text{Hg}) < 2.1 \times 10^{-28} \text{ e cm}. \] (21)

Unfortunately, the implications of the result (21) are somewhat less impressive, due to the electrostatic screening of the nuclear EDM in this essentially Coulomb system. The point is that in a stationary state of such a system, the total electric field acting on each particle must vanish. Thus, an internal rearrangement of the system’s constituents gives rise to an internal field $E_{\text{int}}$ that exactly cancels $E_{\text{ext}}$ at each charged particle; the external field is effectively switched off, and an EDM feels nothing [3, 37, 38].

As to heavy paramagnetic atoms, due to magnetic interactions increasing rapidly with $Z$, the electron EDM therein is not suppressed, but enhanced as $Z^3 \alpha^2$ [39]. In particular, in Tl where the limit (16) was obtained, this enhancement factor reaches $-585$ [40].

However, in the present case of mercury, observable $T$ odd effects are due to the finite size of the nucleus which is small on the atomic scale. To explain how these effects arise here, let us expand the nuclear charge density in powers of the proton coordinates:

\[
\rho(\mathbf{r}) = e \sum_p \delta(\mathbf{r} - \mathbf{r}_p) = e \sum_p \delta(\mathbf{r}) - e \sum_p r_i^p \nabla_i \delta(\mathbf{r})
\]

\[
+ \frac{1}{2} e \sum_p r_i^p r_j^p \nabla_i \nabla_j \delta(\mathbf{r}) - \frac{1}{6} e \sum_p r_i^p r_j^p r_k^p \nabla_i \nabla_j \nabla_k \delta(\mathbf{r}) + \ldots
\] (22)

Going over now to the operators of quadrupole and octupole moments,

\[
Q_{ij} = e \sum_p r_i^p r_j^p - \frac{1}{3} \delta_{ij} e \sum_p r_p^2,
\]

\[
O_{ijk} = e \sum_p \left[ r_i^p r_j^p r_k^p - \frac{1}{5} (r_i^p \delta_{jk} + r_j^p \delta_{ik} + r_k^p \delta_{ij}) r_p^2 \right],
\]
we rewrite formula (22) as

\[ \rho(r) = Z e \delta(r) + \frac{1}{6} Z e (r^2) \Delta \delta(r) - \frac{1}{10} e \sum_p r_p^2 \nabla_i \Delta \delta(r) \]

\[ -d_i \nabla_i \delta(r) + \frac{1}{2} Q_{ij} \nabla_i \nabla_j \delta(r) - \frac{1}{6} O_{ijk} \nabla_i \nabla_j \nabla_k \delta(r) + \ldots; \]  

(23)

in this expression \( \langle r^2 \rangle \) is the mean squared charge radius of the nucleus.

The term \(-e \sum_p r_p^2 \nabla_i \delta(r)\) in the charge density generates P odd, T odd correction to the electrostatic potential. To take into account the mentioned electrostatic screening, we have to subtract from this correction the term

\[ (\mathbf{d} \cdot \nabla) \int \frac{\rho(r)}{|\mathbf{R} - \mathbf{r}|} d^3 \mathbf{r}, \]

where \( \mathbf{d} \) is the EDM of the nucleus. In this way we obtain the following P odd, T odd term in the electrostatic potential:

\[ \phi_{PT}(\mathbf{R}) = -4\pi (\mathbf{S} \cdot \nabla) \delta(\mathbf{R}). \]  

(24)

Here \( \mathbf{S} \) is the so-called Schiff moment

\[ \mathbf{S} = \frac{1}{10} e \sum_p r_p \left( r_p^2 - \frac{5}{3} \langle r^2 \rangle \right). \]  

(25)

The contribution to the Schiff moment from the electric dipole moments \( d_N \) of nucleons can be found as follows. Equation (25) is valid for any system of point-like charges. Let us split the sum in it into the sum over coordinates \( \mathbf{r}_N \) of the centres of mass of nucleons and the sum over coordinates \( \mathbf{\rho}_i \) of point-like charged constituents (say, quarks) of the nucleons:

\[ \mathbf{S} = \frac{1}{10} \sum_N \sum_i e_i \left( (\mathbf{r}_N + \mathbf{\rho}_i)^2 - \frac{5}{3} \langle r^2 \rangle \right) (\mathbf{r}_N + \mathbf{\rho}_i). \]  

(26)

Combining terms of the zeroth and first order in \( \rho \) and taking into account that \( \sum_i e_i = e_N \), \( \sum_i e_i \mathbf{\rho}_i = \mathbf{d}_N \), we present the arising expression for the Schiff moment as a sum of two contributions. The first one is similar to (25), and originates from the P odd, T odd nucleon-nucleon interaction:

\[ \mathbf{S}_1 = \frac{1}{10} e_N \mathbf{r}_N \left( r_N^2 - \frac{5}{3} \langle r^2 \rangle \right); \]  

(27)

here \( e_N = |e| \) for the proton and vanishes for the neutron. The second contribution is due to the internal dipole moments of the nucleons

\[ \mathbf{S}_2 = \frac{1}{6} \sum_N \mathbf{d}_N \left( r_N^2 - \langle r^2 \rangle \right) + \frac{1}{5} \sum_N (\mathbf{r}_N (\mathbf{r}_N \cdot \mathbf{d}_N) - \mathbf{d}_N r_N^2/3). \]  

(28)
If one ascribes the atomic dipole moment to the EDM of the valence neutron in the even-odd nucleus $^{199}$Hg, the corresponding upper limit on the neutron EDM is

$$d_n < 4 \times 10^{-25} \text{ e cm.}$$

It is few times worse than the direct one\(^\text{[41]}\).

On the other hand, though the $^{199}$Hg nucleus in the simple-minded shell model has an odd neutron above a spherically symmetric core, the account for many-body effects allows one to derive an upper limit on the proton EDM

$$d_p < 5.4 \times 10^{-24} \text{ e cm.}$$\(^\text{(29)}\)

This is about an order of magnitude better than the limit obtained from the experiment with TlF molecule\(^\text{[42]}\).

It has been demonstrated, however, that the dipole moments of nuclei induced by the T and P odd nuclear forces can be about two orders of magnitude larger than the dipole moment of an individual nucleon\(^\text{[43]}\). In the simplest approximation of the shell model, where the nuclear spin coincides with the total angular momentum of an odd valence nucleon, while the other nucleons form a spherically symmetric core with the zero angular momentum, the effective T and P odd single-particle potential for the outer nucleon is

$$W = \frac{G}{\sqrt{2}} \frac{\xi}{2m_p} \sigma \cdot \nabla \rho(r).$$\(^\text{(30)}\)

Here \(\xi\) is a dimensionless constant characterizing the strength of the interaction in units of the Fermi weak interaction constant \(G\); \(m_p\), \(\sigma\), and \(r\) are the mass, spin and coordinate of the valence nucleon, respectively.

A simple closed form for the nuclear EDM induced by interaction\(^\text{(30)}\) can be derived as follows. Since the profiles of the nuclear core density \(\rho(r)\) and the potential \(U(r)\) for the outer nucleon are similar, let us assume that they coincide exactly:

$$\rho(r) = U(r) \frac{\rho_0}{U_0}.\text{ (33)}$$

Then the perturbation\(^\text{(30)}\) can be rewritten as

$$W(r) = \lambda \sigma \cdot \nabla U(r), \quad \lambda = \frac{G}{\sqrt{2}} \frac{\xi}{2m_p} \frac{\rho_0}{U_0} = -2 \times 17^{-21} \xi \text{ cm.}$$\(^\text{(31)}\)

Accordingly, the total potential in which the nucleon moves is

$$\tilde{U}(r) = U(r) + W(r) = U(r) + \lambda \sigma \cdot \nabla U(r) = U(|r + \lambda \sigma|).$$\(^\text{(32)}\)

In this potential, it is obvious that the wave function of the external nucleon becomes

$$\tilde{\psi}(r) = \psi(r + \lambda \sigma) = (1 + \lambda \sigma \cdot \nabla) \psi(r),$$\(^\text{(33)}\)
where $\psi(r)$ is its unperturbed value. It is a simple problem now to obtain the following closed expression for the nuclear EDM:

$$d_N = -e(q - Z/A)\lambda \frac{1/2 - K}{I + 1}. \quad (34)$$

Here $Z$, $A$, and $I$ are the nuclear charge, atomic number, and spin, respectively; $K = (l - I)(2I + 1)$, where $l$ is the orbital angular momentum of the outer nucleon; $q = 1$ or 0 for an external proton or neutron, respectively. It is curious that this result of a simple-minded analytical approach is very close to that of the best numerical single-particle calculations.

The characteristic value of the thus induced nuclear EDM is

$$d_N \sim 10^{-21} \xi \text{ cm.} \quad (35)$$

Being interpreted in terms of the CP odd nuclear forces, the experimental result \[21\] leads to the following upper limit:

$$\xi < 0.5 \times 10^{-3}. \quad (36)$$

There are good reasons to assume that the exchange by $\pi^0$-meson is the most efficient mechanism of generating CP odd nuclear forces. This is due to the large value of the strong $\pi NN$ coupling constant $g_s = 13.5$ and to the small pion mass, as well as to the fact that outer proton and neutron orbitals in heavy nuclei are quite different. The P odd, T odd effective $\pi NN$ Lagrangians are conveniently classified by their isotopic properties \[44, 45, 46\]:

$$\Delta T = 0. \quad L_0 = \bar{g}_0 \left[ \sqrt{2} (\bar{p}n\pi^+ + \bar{p}p\pi^-) + (\bar{p}p - \bar{p}n) \pi^0 \right]; \quad (37)$$

$$|\Delta T| = 1. \quad L_1 = \bar{g}_1 \bar{N}N \pi^0 = \bar{g}_1 (\bar{p}p + \bar{m}n) \pi^0; \quad (38)$$

$$|\Delta T| = 2. \quad L_2 = \bar{g}_2 \left[ \sqrt{2} (\bar{p}m\pi^+ + \bar{m}p\pi^-) - 2(\bar{p}p - \bar{m}n) \pi^0 \right]. \quad (39)$$

With the strong interaction Lagrangian given by formula \[11\], it can be easily seen that in heavy nuclei, at least in the single-particle approximation, the contributions of the P odd, T odd interactions $L_0$, $L_2$, which conserve isospin and change it by 2, are proportional to $N - Z$ and thus are suppressed as compared to $L_1$ interaction changing isospin by 1, which is proportional to $N + Z = A$. For the present case of the $^{199}$Hg nucleus there is an additional occasional cancellation in the isoscalar contribution. In this way the limit \[36\] can be transformed to

$$g_1 < 0.5 \times 10^{-11}. \quad (40)$$

The Standard Model (SM) prediction for this constant is \[47\]

$$g_1 \sim 10^{-17}. \quad (41)$$

Thus, the theoretical predictions of the SM for dipole moments and CP odd nuclear forces are about six orders of magnitude below the present experimental upper limits on them.
In fact, this means that the searches for electric dipole moments now, at the present level of accuracy, are extremely sensitive to possible new physics.

Theoretical models of CP violation are too numerous to discuss all of them, and most of them have too many degrees of freedom. It is convenient therefore to proceed in a phenomenological way: to construct CP odd quark-quark, quark-gluon, and gluon-gluon operators of low dimension, and find upper limits on the corresponding coupling constants from the experimental results for \( d_n \) and \( d^{(199)\text{Hg}} \). The analysis performed in [47, 48], has demonstrated that the limits on the effective CP odd interaction operators obtained from the neutron and atomic experiments are quite comparable. These limits are very impressive. All the constants are several orders of magnitude less than the usual Fermi weak interaction constant \( G \). In particular, these limits strongly constrain some popular models of CP violation, such as the model of spontaneous CP violation in the Higgs sector, and the model of CP violation in the supersymmetric \( S0(10) \) model of grand unification.

4 Electric dipole moments, storage rings and linear traps

The various upper limits on EDMs set so far, constitute a valuable contribution to elementary particle physics and to our knowledge of how the Nature is arranged; the null results obtained so far are important. But it is only natural to think of essential progress in the field, of finding a positive result, of eventually discovering permanent electric dipole moment. In particular, it would be tempting to get rid of the electrostatic screening of nuclear EDMs. So, let me add to the above stories, a new one. It should be started with the discussion of

4.1 Idea of new muon EDM experiment

The idea is to search for the muon EDM in a storage ring, with muons in it having natural longitudinal polarization [35]. An additional spin precession due to the EDM interaction with external field should be monitored by counting the decay electrons, their momenta being correlated with the muon spin, due to parity nonconservation in the muon decay. The frequency \( \omega \) of the spin precession with respect to the particle momentum in external magnetic and electric fields, \( \mathbf{B} \) and \( \mathbf{E} \), is

\[
\omega = -\frac{e}{m} \left[ a\mathbf{B} - a \frac{\gamma}{\gamma + 1} \mathbf{v} \cdot \mathbf{B} - \left( a - \frac{1}{\gamma^2 - 1} \right) \mathbf{v} \times \mathbf{E} \right] - \eta \frac{e}{m} \left[ \mathbf{E} - \frac{\gamma}{\gamma + 1} \mathbf{v} \cdot \mathbf{E} + \mathbf{v} \times \mathbf{B} \right].
\]

Here the anomalous magnetic moment \( a \) is related to the \( g \)-factor as follows: \( a = g/2 - 1 \) (for muon \( a = \alpha/2\pi \)); \( \mathbf{v} \) is the particle velocity; \( \gamma = 1/\sqrt{1 - v^2} \). The last line in this
formula describes the precession due to the EDM $d$, the dimensionless constant $\eta$ being related to $d$ as follows:

$$d = \frac{e}{2m} \eta.$$

This last line can be obtained from the terms proportional to the anomalous magnetic moment in (42) by substituting $a \rightarrow \eta$ and changing to dual fields: $B \rightarrow E, \ E \rightarrow -B$. Expression (42) simplifies in the obvious way for $(\mathbf{v} \cdot B) = (\mathbf{v} \cdot E) = 0$. Just this case is considered below.

The idea of [35] is to compensate for the usual spin precession in the vertical magnetic field $B$ by the precession in a radial electric field $E$, i.e. to choose $E$ in such a way that the first line in (42) vanishes at all. Then the spin precession with respect to momentum is due only to the EDM interaction with the vertical magnetic field, and since electric fields in a storage ring are much smaller than magnetic ones, it reduces to

$$\omega = \omega_e = -\frac{e}{m} \eta \mathbf{v} \times \mathbf{B}. \quad (43)$$

In a nutshell, due to the EDM interaction, the muon spin precesses around the motional electric field $\mathbf{v} \times \mathbf{B}$. In this way the spin acquires a vertical component which linearly grows with time. The P odd correlation of the decay electron momentum with the muon spin leads to the difference between the number of electrons registered above and below the orbit plane.

In [35], it is stated that the sensitivity to the muon EDM can be improved in the planned experiment by six orders of magnitude, to the level of $10^{-24} \, e \, cm$.

But after all, the useful signal here is due to the spin precession in the electric field in the muon rest frame. So, the question is: what about the screening of the electric field which is responsible in particular for reducing the record-breaking result (21) to much more modest bound on $d_n$? The explanation is as follows. All previous EDM experiments consisted essentially in the measurements of frequencies, i.e. referred to stationary states. In the present case the spin precession itself is being measured, i.e. this is completely nonstationary situation. Thus there is no screening suppression here.

### 4.2 Nuclear electric dipole moments at ion storage rings

In the same way one can search for an EDM of a polarized $\beta$-active nucleus in a storage ring [49]. In this case as well, the precession of nuclear spin due to the EDM interaction can be monitored by the direction of the $\beta$-electron momentum. $\beta$-active nuclei have serious advantages as compared to muon. The life time of a $\beta$-active nucleus can exceed by many orders of magnitude that of a muon. The characteristic depolarization time of the ion beam can reach few seconds, and is also much larger than the muon life time, which is about $10^{-6} \, s$. Correspondingly, the angle of the rotation of nuclear spin, which is due to the EDM interaction and which accumulates with time, may be also by orders of magnitude larger than that of a muon. By the same reason of the larger life time, the quality of an ion beam can be made much better than that of a muon beam.
\[
I^+ \to I'^\pm \quad \mu \quad z \quad a \times 10^3 \quad t_{1/2} \quad Q(\text{barn}) \quad \text{branching}
\]

| \( ^{131}_{\text{I}} \) | \( 7/2^+ \rightarrow 5/2^+ \) | 2.742(1) | 51 | -1.9(0.4) | 8.0 d | -0.40 | 90% |
| \( ^{133}_{\text{I}} \) | \( 7/2^+ \rightarrow 5/2^+ \) | 2.856(5) | 53 | 16(2) | 21 h | -0.27 | 83% |
| \( ^{139}_{\text{Cs}} \) | \( 7/2^+ \rightarrow 7/2^- \) | 2.696(4) | 53 | 2(1) | 9.3 m | -0.075 | 82% |
| \( ^{223}_{\text{Fr}} \) | \( 3/2^- \rightarrow 3/2^- \) | 1.17(2) | 87 | -7\pm20 | 22 m | 1.2 | 67% |

Table 2

However, necessary conditions here are also quite serious.

First of all, to make realistic the mentioned compensation of the EDM-independent spin precession by a relatively small electric field, the effective nuclear g-factor should be close to 2 (as this is the case for the muon). For a nucleus with the total charge \( Ze \), mass \( Am_p \), spin \( I \), and magnetic moment \( \mu \), the effective anomalous magnetic moment is

\[
a = \frac{g}{2} - 1 = \frac{A}{2} \frac{\mu}{I} - 1.
\]

Some fine-tuning of \( a \) is possible sometimes by taking, instead of a bare nucleus, an ion with closed electron shells. An accurate formula for the anomaly of an ion with the total charge \( z \), is

\[
a = \frac{A}{2z} \frac{\mu}{I} - 1.00722 + \frac{\Delta}{Am_p} - \frac{z m_e}{A m_p}.
\]

It includes the correction for the atomic mass excess \( \Delta \).

It is more practical perhaps to confine, in line with bare nuclei, to helium-like ions, i.e. to \( z = Z, \ Z - 2 \). One of the reasons is the same electrostatic screening. It is complete for a neutral atom, and for an ion it is only partial, being proportional to the number of electrons \( Z - z \). So, the smaller this number is, the better for our problem.

Some ions which look at the moment promising from the point of view of the EDM searches are presented in the Table. More complete list, comprising about 30 candidates, can be found in [49].

The errors in the values of anomalous magnetic moments \( a \), presented in the Table, correspond to the experimental errors in values of \( \mu \). The \( \beta^- \) branchings are indicated in the last column.

Let us note that a background due to nuclear quadrupole moments, even as large as \( Q \sim 1 \) barn, is not dangerous at reasonable parameters of a storage ring even for the EDM sensitivity as high as \( 10^{-26} \) e cm.

It looks at the moment that the most serious problem for both muon and nuclear EDM experiments consists in extremely strict demands on the alignment of the radial electric field. According to I. Koop and Y. Semertzidis, its spurious vertical component is bounded as follows:

\[
\frac{E_z}{E_r} < \eta \frac{v}{a}.
\]
For the muon experiment this ratio should not exceed $10^{-8}$. Such an accuracy is difficult to attain. The situation is better for ions at the same ratio $v/a$. The point is that with the same EDM value, $\eta$ for a nucleus is larger by a factor $(A/z)(m_p/m_\mu) \sim 20$. However, even for ions the problem is very serious.

An excellent idea by I. Koop is to work here without electric field at all. As to muons, their life time is short, $2 \times 10^{-6}$ s. So, during the measurement their spins anyway will make only few turns due to $a$. Therefore, in this case it looks reasonable to make the measurement time smaller, i. e. to sacrifice an order of magnitude in sensitivity, but to get rid of the formidable problem. As to ions, the idea works differently. Here the residual spin precession due to $a$ can be used for modulation of the useful signal, i. e. of the spin precession due to the EDM. The modulation of the signal partly compensates for the loss in its magnitude.

Now, if a sufficiently large EDM signal can be attained, i. e. if the angle of the spin rotation can reach, say, a milliradian, one could think about an experiment with stable nuclei or nuclei of a large life time. Their polarization could be measured in scattering experiments (the idea advocated by Y. Semertzidis and A. Skrinsky). Here the deuteron is a prominent candidate [50] in spite of its not so small anomaly $a = -0.143$. The deuteron fluxes are attainable exceeding $10^{12}$ particles per second, with high degree of polarization. The deuteron polarization is measured now with an accuracy $\sim 10^{-2}$. From the theoretical point of view, the calculation of its EDM $d_d$ is a relatively clean problem. If induced by P odd, T odd nuclear forces, $d_d$ can be calculated in the same way as the deuteron AM with the following result [22]:

$$d_d = -\frac{eg_1}{12\pi m_\pi} \frac{1 + \xi}{(1 + 2\xi)^2} I.$$ \hspace{1cm} (46)

Numerically, it is

$$d_d = -2.4 \times 10^{-14} g_1 e \text{ cm.}$$ \hspace{1cm} (47)

Calculations with more realistic (than zero-range approximation) deuteron wave functions allows one to estimate the accuracy of this result as 30%. Thus, the deuteron EDM is due essentially to the same P odd, T odd $\pi NN$ constant $g_1$ as the EDM of a heavy nucleus.

On the other hand, the deuteron EDM can arise due to the proton and neutron dipole moments. With the deuteron being essentially a $^3S_1$ bound state, this contribution to its EDM is

$$d_d(n, p) = d_p + d_n.$$ \hspace{1cm} (48)

If one assumes that the nucleon dipole moments are also due to the P odd, T odd $\pi NN$ interaction, they can be calculated in the chiral limit as well [51]. The corresponding results, written in terms of the constants introduced in formulae (11), (37), (39), are

$$d_n = -d_p = \frac{e}{m} \frac{g(g_0 + g_2)}{4\pi^2} \ln \frac{m}{m_\pi}.$$ \hspace{1cm} (49)

The idea to use a dedicated deuteron storage ring for measuring $d_d$ with the sensitivity of $10^{-26} e \text{ cm}$ is now seriously discussed by experimentalists (see [http://www2.bnl.gov/~muonedm/]).
This experiment would be a big leap forward in the investigation of the nature of CP violation.

On the one hand, it would give immediately the information on the sum of the EDMs \( d_p + d_n \) on the same level of \( 10^{-26} \text{ cm} \), a gain by an order of magnitude as compared to (15), (29).

On the other hand, being interpreted in terms of T and P odd nuclear forces, it would correspond in virtue of (47) to the sensitivity for \( g_1 \) an order of magnitude better than the result (40) of the mercury experiment.

The great importance of the deuteron EDM experiment is obvious.

And at last

### 4.3 Linear electrostatic trap for heavy polar molecules

We have mentioned already the formidable problem of the electric field alignment (see (45)), perhaps the most serious one for the storage ring EDM experiments. A cardinal solution of the problem of spurious spin rotation would be to get rid of external magnetic fields at all and to make the velocity as small as possible to suppress the motional magnetic field in the rest frame. Both these aims are reached, at least in principle, in the proposal to use linear electrostatic trap for cold polar molecules [52].

The trap looks as a long capacitor where cold molecules move along \( z \) axis. The electric field \( E \) is constant on the \( z \) axis and directed along \( y \), it changes only at the ends, thus forming a longitudinal potential well. On the other hand, the profile of \( E \) changes in the \( xy \) plane in such a way as to guarantee a sort of strong focusing for molecules aligned along \( y \) and moving along \( z \) axis. Due to the essentially one-dimensional motion of the molecules, one also gets rid (at least, to first approximation) of one more serious problem, that of the inertial dragging of spin.

The nuclear spin (or electron one in a radical) precesses in the \( xz \) plane due to the EDM interaction with the effective electric field. It is not the capacitor field \( E_y \), but the intermolecular one which is huge, it can amount to \( 10^9 \text{ V/cm} \). This last advantage of experiments with polar molecules was in fact pointed out long ago [53].

Obviously, the potential advantages of this idea are huge, and it well deserves serious experimental investigations.

### References

[1] Ya.B. Zel’dovich, Zh. Eksp. Teor. Fiz. 33 (1957) 1531 [Sov. Phys. JETP 6 (1957) 1184] (the paper contains also the mention of the analogous results obtained by V.G. Vaks.).

[2] C.S. Wood et al., Science 275 (1997) 1759.

[3] E.M. Purcell and N.F. Ramsey, Phys. Rev., 78 (1950) 807.
[4] V.V. Flambaum and I.B. Khriplovich, Zh. Eksp. Teor. Fiz. 79 (1980) 1656
[Sov. Phys. JETP 52 (1980) 835].

[5] V.V. Flambaum, I.B. Khriplovich and O.P. Sushkov, Phys. Lett. B146 (1984) 367.

[6] P.A. Frantsuzov and I.B. Khriplovich, Z. Phys. D7 (1988) 297.

[7] A.Ya. Kraftmakher, Phys. Lett. A132 (1988) 167.

[8] B. Desplanques, J.F. Donoghue and B.R. Holstein, Ann. Phys. 124 (1980) 449.

[9] V.F. Dmitriev et al., Phys. Lett. B125 (1983) 1.

[10] V.V. Flambaum, V.B. Telitsin and O.P. Sushkov, Nucl. Phys. A444 (1985) 611.

[11] V.F. Dmitriev and I.B. Khriplovich, nucl-th/0201041

[12] V.F. Dmitriev, I.B. Khriplovich and V.B. Telitsin, Nucl. Phys. A577 (1994) 691.

[13] N. Auerbach and B.A. Brown, Phys. Rev. C60 (1999) 025501.

[14] V.F. Dmitriev and V.B. Telitsin, Nucl. Phys. A674 (2000) 168.

[15] W.C. Haxton, C.-P. Liu and M.J. Ramsey-Musolf, Phys. Rev. C65 (2002) 045502.

[16] L. Rydstrom et al., Nucl. Phys. A512 (1990) 217.

[17] E.G. Adelberger, W.C. Haxton, Ann.Rev.Nucl.Part.Sci. 35 (1985) 501.

[18] P.A. Vetter et al.,Phys. Rev. Lett. 74 (1995) 2658.

[19] M.S. Hussein, A.F.R. de Toledo Piza, O.K. Vorov, A.K. Kerman, Phys. Rev. C60
(1999) 064615.

[20] M.S. Hussein, A.F.R. de Toledo Piza, O.K. Vorov, A.K.Kerman, Nucl. Phys. A686
(2001) 163.

[21] K.E. Arinstein, V.F. Dmitriev, I.B. Khriplovich, Yad. Fiz., in press.

[22] I.B. Khriplovich and R.V. Korkin, Nucl. Phys. A665 (2000) 365.

[23] M.J. Savage and R.P. Springer, Nucl. Phys. A644 (1998) 235;
erratum Nucl. Phys. A657 (1999) 457.

[24] D. DeMille, Phys. Rev. Lett. 74 (1995) 4165.

[25] S. Porsev, Yu. Rakhлина, and M.Kozlov, Pis’ma Zh. Eksp. Teor. Fiz. 61 (1995) 449
[Sov. Phys. JETP Pis’ma Zh 61 (1995) 459].

[26] B.P. Das, Phys. Rev. A56 (1997) 1635.
[27] C.J. Bowers et al., Phys.Rev. A53 (1996) 3103.
[28] C.J. Bowers et al., Phys.Rev. A59 (1999) 3513.
[29] J.E. Stalnaker et al., Phys. Rev. A65 (2002) 031403.
[30] K.F. Smith et al., Phys. Lett. B234 (1990) 191.
[31] I.S. Altarev et al., Phys. Lett. B276 (1992) 242.
[32] P.G. Harris et al., Phys. Rev. Lett. 82 (1999) 904.
[33] B.C. Regan et al., Phys. Rev. Lett. 88 (2002) 0071805.
[34] J. Baily et al., Zs. Phys. G4 (1978) 345.
[35] Y.K. Semertzidis, in Proceedings of the Workshop on Frontier Tests of Quantum Electrodynamics and Physics of the Vacuum (1998, Sandansky, Bulgaria).
[36] M.V. Romalis, W.C. Griffith and E.N. Fortson, Phys. Rev. Lett. 86 (2001) 2505.
[37] R.L. Garwin and L.M. Lederman, Nuovo Cim. 11 (1959) 776.
[38] L.I. Schiff, Phys. Rev. 132 (1963) 2194.
[39] P.G.H. Sandars, Phys. Lett. 14 (1965) 194.
[40] Z.W. Liu and H.P. Kelly, Phys. Rev. A45 (1992) R4210.
[41] V.F. Dmitriev and R.A. Sen’kov, "nucl-th/0306050." 
[42] D. Cho, K. Sangster and E. Hinds, Phys. Rev. Lett. 63 (1989) 2559.
[43] O.P. Sushkov, V.V. Flambaum and I.B. Khriplovich, Zh. Eksp. Teor. Fiz. 87 (1984) 1521; [Sov. Phys. JETP, 60 (1984) 873].
[44] G. Barton, Nuovo Cim., 19 (1961) 512.
[45] W. C. Haxton and E. M. Henley, Phys. Rev. Lett., 51 (1983) 1937.
[46] P. Herczeg, in Tests of Time Reversal Invariance in Neutron Physics, edited by N. R. Robertson, C. R. Gould, and J. D. Bowman (World Scientific, Singapore, 1987), p. 24; see also P. Herczeg, Hyperfine Interactions 43 (1988) 77.
[47] V.M. Khatsymovsky, I.B. Khriplovich and A.S. Yelkhovsky, Ann. Phys. 186 (1988) 1.
[48] V.M. Khatsymovsky and I.B. Khriplovich, Phys. Lett. B296 (1992) 219.
[49] I.B. Khriplovich, Phys. Lett. B444 (1998) 98; Nucl. Phys. A663&664 (2000) 147c.
[50] I.B. Khriplovich, in Proceedings of the NATO Advanced Study Institute on Trapped Particles and Fundamental Physics, Les Houches, France, May 23-June 2, 2000 (2002 Kluwer Academic Publishers).

[51] R.J. Crewther, P. Di Vecchia, G. Veneziano, and E. Witten, Phys. Lett. B88 (1979) 123; 91 (1980) 487(E).

[52] I.A. Koop, to be published.

[53] P.G.H. Sandars, Phys. Rev. Lett. 19 (1967) 1396.