Effects of core polarization on the nuclear Schiff moment

V.F. Dmitriev, R.A. Sen'kov

Budker Institute of Nuclear Physics, Novosibirsk-90, 630090 Russia and Novosibirsk State University, Novosibirsk-90, 630090 Russia

N. Auerbach

School of Physics and Astronomy, Tel Aviv University, Tel-Aviv, 69978 Israel

Schiff moments were calculated for a set of nuclei with full account of core polarization effects. A finite range P and T violating weak nucleon-nucleon interaction has been used in the calculations. While in the absence of core polarization the Schiff moment depends on one combination of the weak interaction constants, in the presence of core polarization the Schiff moment depends on all three constants separately. The dominant contribution comes from isovector, $\Delta T = 1$, part of the weak interaction. The effects of core polarization were found to have in general a large effect on the Schiff moments.

PACS numbers: 21.10.Ky, 21.60.Iz

I. INTRODUCTION

The theory of static atomic dipole moments relies mainly on presence of parity (P) and time (T) reversal violating electro-magnetic (E-M) moments generated by nuclei. It is conjectured that the dominant contribution to the atomic dipole moment comes from parity and time reversal violating components in the nucleon-nucleon (N-N) force. These components of the interaction induce P-T violating E-M moments in the nucleus. In turn these forbidden nuclear moments produce parity and time reversal violation in the E-M field which can induce static electric dipole moments in the atom. This way of reasoning is the working assumption in the study of the atomic dipole moment [1, 2].

The measurements of the atomic dipole moment are presently very advanced and reached a very high degree of precision [1, 3]. These experiments are performed with neutral atoms. It has been shown years ago [4] that in such cases the field of the electric dipole moment of the nucleus is screened by the field of the electrons and there is no effect when the atom is placed in an external electric field. This situation applies strictly only when one considers point-like particles. When the finite size of the nucleus is taken into account one can go to higher orders in the expansion of the E-M field and find that in the next order in the expansion it is the Schiff moment [5] that can induce an electric dipole moment in the atom. (See next section).

The upper limits for the existence of dipole moments obtained in the atomic measurements cannot provide directly limits on the presence of T and P violation in the nuclear Hamiltonian. One needs theoretical input. In order to make the connection one must have precise calculations of moments in nuclei produced by the T-P non-conserving part of the N-N interaction.

In the past the theory of Schiff moments in spherical nuclei was limited to either a single-particle value (in Z-odd, N-even nuclei) or calculated for a simple two particle – one hole (2p-1h) configuration (in Z-even, N-odd nuclei).

In this paper we will go beyond these simple calculations and include various effects of core polarization. As we will see these effects are large compared to the results obtained in lowest order.

We should stress that an important aspect of this work is the use of a finite range interaction for the P-T-violating component of the N-N force. A finite range interaction was used recently in the calculation of the Schiff moment of ^{199}Hg [6] and ^{225}Ra[7]. In the present work we examine the influence of finite range for a number of nuclei and discuss more general features that this leads to.

We compute the Schiff moments for the following nuclei, for reasons explained below.

A. Neutron-odd nuclei

a. ^{199}*Hg:* The most accurate upper limit on parity and time reversal violation has been obtained from the measurement of the atomic electric dipole moment in the Hg atom [3]. The nucleus has (Z=80) has an odd number of neutrons. The lowest order contribution to the Schiff moment comes from a 2p-1h configuration (a proton 1p-1h added to the neutron) admixed into the ground state (g.s.). Because of the importance of this nucleus in the present experimental studies we have examined in more detail the influence of core polarization effects on the Schiff moment .

b. ^{129}Xe: The A=129 isotope of Xenon is also an even-Z, odd-N nucleus and the lowest order contribution to the Schiff moment comes from a 1p-1h proton excitation of the core. Similarly to ^{199}Hg the atom of ^{129}Xe is used in experimental studies of the atomic electric dipole moment [1]. The upper limit for the dipole moment obtained in the experiment is only several times higher than in Hg. We have, therefore as in the case of Hg, found it useful to calculate more accurately the Schiff moment by introducing several effects coming from core polarization and not included in previous theoretical studies.

c. ^{211}Rn and ^{213}Ra: The light isotopes of radon and radium are considered to be transitional nuclei, still being close to spherical in their ground states [8, 9]. The heavier isotopes of these elements (A=221-225) are found to have quadrupole plus octupole deformed shapes in the ground states [8, 9]. It was suggested that in these heavier isotopes of Rn and Ra (as well as some other elements in this region) the Schiff moments are strongly enhanced because of the quadrupole+octupole deformations [10, 11]. It is therefore useful for the purpose of comparison to calculate accurately the Schiff moments in the neighboring, lighter isotopes where the onset of deformations has not yet occurred.

d. ^{225}Ra: Here we calculate the Schiff moment assuming that the nucleus is spherical. We do this in order to have a reference value so that one can see the degree of enhancement when a full calculation , that includes deformations, is performed.

B. Proton-odd nuclei

In a Z-odd, N-even nucleus the odd proton can contribute directly to the Schiff moment. It is therefore usually assumed that the dominant contribution is the single-particle value of the last proton.

e. ^{133}Cs: We examined this point for ^{133}Cs (Z=55) where the core-polarization correction was evaluated together with the single-particle Schiff moment.

f. ^{223}Fr: This nucleus (Z=87) was found to have an enhanced Schiff moment when the quadrupole+octupole deformations were taken into account [11]. As in the case of ^{225}Ra for the same reason we computed the single-particle and core polarization values taking ^{223}Fr to be spherical.

II. P AND T VIOLATING NUCLEON-NUCLEON INTERACTION

The expectation value of the Schiff moment operator [5]

$$\hat{\mathbf{S}} = \frac{1}{10}\sum_i e_i(r_i^2\mathbf{r}_i - \frac{5}{3} < r^2 >_{ch} \mathbf{r}_i) \tag{1}$$

can be non zero only in the presence of P and T violating nuclear forces. There are good reasons to assume that the exchange of a π^0-meson is the most efficient mechanism of generating CP odd nuclear forces. This is due to the large value of the strong πNN 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 πNN Lagrangians are conveniently classified by their isotopic properties [12–14]:

$$\Delta T = 0. \quad L_0 \;=\; g_0\,[\,\sqrt{2}\,(\overline{p}n\,\pi^+ + \overline{n}p\,\pi^-)\;+(\,\overline{p}p - \overline{n}n)\,\pi^0]; \tag{2}$$

$$|\Delta T| = 1. \quad L_1 \;=\; g_1\,\overline{N}N\,\pi^0 \;=\; g_1\,(\,\overline{p}p\;+\overline{n}n)\,\pi^0; \tag{3}$$

$$|\Delta T| = 2. \quad L_2 \;=\; g_2\,[\,\sqrt{2}\,(\overline{p}n\,\pi^+ + \overline{n}p\,\pi^-)\;-2(\,\overline{p}p - \overline{n}n)\,\pi^0]. \tag{4}$$

The above Lagrangians generate a P and T odd nucleon-nucleon interaction of the form

$$\begin{aligned} W(\mathbf{r}_a - \mathbf{r}_b) = -\frac{g_s}{8\pi m_p}\,&[(g_0\boldsymbol{\tau}_a\cdot\boldsymbol{\tau}_b + g_2(\boldsymbol{\tau}_a\cdot\boldsymbol{\tau}_b - 3\tau_a^z\tau_b^z))(\boldsymbol{\sigma}_a - \boldsymbol{\sigma}_b) \\ &+g_1(\tau_a^z\boldsymbol{\sigma}_a - \tau_b^z\boldsymbol{\sigma}_b)]\cdot\boldsymbol{\nabla}_a\frac{e^{-m_\pi r_{ab}}}{r_{ab}}, \end{aligned} \tag{5}$$

where m_p is the proton mass and $r_{ab} = |\mathbf{r}_a - \mathbf{r}_b|$.

Together with the finite range interaction Eq.(5) a phenomenological zero range effective interaction has often been used in estimates of the Schiff moment. It has the form [15]

$$\begin{aligned} W_c(\mathbf{r}_a - \mathbf{r}_b) = \frac{G}{\sqrt{2}}\frac{1}{2m_p}\,&((\eta_{ab}\boldsymbol{\sigma}_a - \eta_{ba}\boldsymbol{\sigma}_b)\cdot\;\boldsymbol{\nabla}_a\delta(\mathbf{r}_a - \mathbf{r}_b)+ \\ &\eta'_{ab}[\boldsymbol{\sigma}_a\times\boldsymbol{\sigma}_b]\cdot\{(\mathbf{p}_a - \mathbf{p}_b),\delta(\mathbf{r}_a - \mathbf{r}_b)\})\,, \end{aligned} \tag{6}$$

where G is the Fermi constant. The interaction (6) can be obtained from (5) in the limit $m_\pi \;\to\; \infty$. In this limit we have the following correspondence between the interaction constants η_{ab} and g_i

$$\eta_{pp} = \frac{\sqrt{2}}{Gm_\pi^2}g_s(-g_0 + 2g_2 - g_1), \quad \eta_{nn} = \frac{\sqrt{2}}{Gm_\pi^2}g_s(-g_0 + 2g_2 + g_1),$$

$$\eta_{np} = \frac{\sqrt{2}}{Gm_\pi^2} g_s(g_0 - 2g_2 + g_1), \quad \eta_{pn} = \frac{\sqrt{2}}{Gm_\pi^2} g_s(g_0 - 2g_2 - g_1). \tag{7}$$

The second term in Eq.(6) comes from the exchange matrix elements of the interaction (5) in the limit $m_\pi \to \infty$

$$\eta'_{pp} = \frac{\sqrt{2}}{Gm_\pi^2} g_s(g_0 + 2g_2 + g_1), \quad \eta'_{nn} = \frac{\sqrt{2}}{Gm_\pi^2} g_s(g_0 + 2g_2 - g_1),$$

$$\eta'_{pn} = \eta'_{np} = \frac{\sqrt{2}}{Gm_\pi^2} 2g_s(g_0 - g_2). \tag{8}$$

One should however remember that the correspondence given by (7) and (8) is exact only in the limit of infinite pion mass. For finite pion mass this correspondence is somewhat conditional since the matrix elements of finite range interaction differ from the matrix elements of zero range interaction. Especially large is the difference between the exchange matrix elements.

For a nucleus having an odd proton there are two kinds of contributions to the nuclear Schiff moment even in the absence of a strong residual interaction between the odd proton and the nucleons in the core [16]. One contribution comes from the odd proton. We call it a single particle contribution. Another contribution comes from the core nucleons.

The weak interaction, Eq.(5), generates a weak correction to the nuclear mean field. The correction is [6]

$$\delta U(\mathbf{r}) = \frac{g_s m_\pi^2}{\pi m_p} (\boldsymbol{\sigma} \cdot \mathbf{n}) \tau^z \int_0^\infty r'^2 dr' \; b(r, r')[(g_0 - 2g_2)(\rho_p(r) - \rho_n(r)) +$$
$$g_1(\rho_p(r) + \rho_n(r))], \tag{9}$$

where $b(r, r')$ is a combination of spherical Bessel functions of an imaginary argument $i_n(x) = \sqrt{\pi/2x} I_{n+1/2}(x)$, and $k_n(x) = \sqrt{2/\pi x} K_{n+1/2}(x)$

$$b(r, r') = i_1(m_\pi r) k_0(m_\pi r') \theta(r' - r) - i_0(m_\pi r') k_1(m_\pi r) \theta(r - r').$$

Note, that the weak potential given by Eq.(9) is proportional to τ^z. It is pure isovector. The contribution of $\Delta T = 0, 2$ channels to the weak potential is proportional to $\rho_p - \rho_n$, while the contribution of $\Delta T = 1$ channel is proportional to $\rho_p + \rho_n$. It means that the weak potential is most sensitive to g_1. The contributions of g_0 and g_2 are suppressed by the factor $(N - Z)/A$. This suppression has an important effect on the $\Delta T = 0$ channel. The finite range real NN-interaction (5) has both direct and exchange matrix elements. The

TABLE I: The contributions from the direct and exchange potentials to the single particle Schiff moment of ^{133}Cs. The units are e fm^3.

	$g_s g_0$	$g_s g_1$	$g_s g_2$
direct	0.011	-0.109	-0.022
exchange	-0.016	-0.003	-0.006

direct matrix elements contribute to the potential in Eq.(9). For a finite range interaction the exchange matrix elements are usually small and in many cases can be safely omitted. For the $\Delta T = 0$ channel, due to suppression in the direct potential, the exchange and direct potentials become comparable.

The weak potential (9) produces a correction $\delta\psi_\nu(\mathbf{r})$ to the single particle wave function $\psi_\nu(\mathbf{r})$ of the odd proton. With this correction the expectation value of the Schiff moment is

$$S_{sp} = \langle \psi_\nu + \delta\psi_\nu | \hat{S}_z | \psi_\nu + \delta\psi_\nu \rangle = \langle \delta\psi_\nu | \hat{S}_z | \psi_\nu \rangle + \langle \psi_\nu | \hat{S}_z | \delta\psi_\nu \rangle. \tag{10}$$

Note that this single particle contribution is absent for neutron odd nuclei. In Table 1 we show the single particle contributions to the Schiff moment of ^{133}Cs from the direct and exchange potentials in all three channels. One can see that while in $\Delta T = 1$ and 2 channels the exchange contribution is really small, in $\Delta T = 0$ channel it is even larger in absolute value than the zero order term and has opposite sign. The same behaviour was found also for ^{223}Fr nucleus.

The weak potential δU is defined by diagonal matrix elements of the two body weak interaction (5) over core states. Another type of contribution to the Schiff moment comes from off diagonal matrix elements of the two body weak interaction.

$$S_{core} = \sum_{1p1h} \frac{\langle \nu, 0 | \hat{S}_z | 1p1h, \nu \rangle \langle \nu, 1p1h | W | 0, \nu \rangle + \langle \nu, 0 | W | 1p1h, \nu \rangle \langle \nu, 1p1h | \hat{S}_z | 0, \nu \rangle}{E_0 - E_{1p1h}}. \tag{11}$$

Calculating the matrix elements in (11) we obtain the following equation for this contribution

$$S_{core} = \sum_{\nu_1 \nu_2} \langle \nu, \nu_1 | W | \nu_2, \nu \rangle \frac{n_{\nu_1} - n_{\nu_2}}{\epsilon_{\nu_1} - \epsilon_{\nu_2}} \langle \nu_2 | \hat{S}_z | \nu_1 \rangle, \tag{12}$$

where n_ν are the occupation numbers and ϵ_ν are the single particle energies.The states ν_1 and ν_2 belong to protons, therefore, for odd protons nuclei S_{core} is proportional to η_{pp} while for

TABLE II: The single particle and the core contributions for two proton odd nuclei. Finite range weak interaction. The units are e fm^3

		$g_s g_0$	$g_s g_1$	$g_s g_2$
^{133}Cs	S_{sp}	-0.005	-0.112	-0.028
^{133}Cs	S_{core}	0.088	0.088	-0.176
^{223}Fr	S_{sp}	0.014	0.084	0.029
^{223}Fr	S_{core}	-0.135	-0.135	0.27

TABLE III: The single particle and the core contributions for two proton odd nuclei. Zero range weak interaction. The units are e fm^3

		$g_s g_0$	$g_s g_1$	$g_s g_2$
^{133}Cs	S_{sp}	-0.06	-0.23	-0.0562
^{133}Cs	S_{core}	0.18	0.18	-0.36
^{223}Fr	S_{sp}	0.071	0.19	0.072
^{223}Fr	S_{core}	-0.209	-0.209	0.418

odd neutrons nuclei it is proportional to η_{np}. The magnitudes of S_{sp} and S_{core} are comparable for proton odd nuclei. This is in contrast to the P odd and T even nuclear anapole moment where the core contribution is smaller by a factor $1/A^{1/3}$. In Table 2 we show the singe particle and the core contributions for two proton odd nuclei. The sum of three terms in the core contribution is proportional to the combination $-g_0 - g_1 + 2g_2 \sim \eta_{pp} = -\eta_{np}$.

At this place it is worth to compare the single-particle and the core contributions calculated in the limit of zero range weak interaction. The results of calculations are presented in Table III. We see that the difference is approximatly by a factor 2. The matrix elements of a finite range interaction are smaller due to smaller overlap of the wave functions in matrix elements of the interaction. However, this is not a general rule. For the Hg nucleus, as we mentioned in [6] the effect of a finite range is insignificant and goes even in opposite direction. For a finite range interaction we obtained

$$S_{core} = -0.085(g_s g_0 + g_s g_1 - 2g_s g_2),$$

while for zero range interaction

$$S_{core} = -0.058(g_s g_0 + g_s g_1 - 2g_s g_2).$$

For Rn nucleus the difference is 38% and for Xe nucleus it becomes again close to 2.

III. CORE POLARIZATION

The above results were obtained in a simple independent particle model (IPM). In our calculations we used a full single-particle spectrum including continuum. The single particle basis was obtained using the partially self-consistent mean-field potential of [19]. The potential includes four terms. The isoscalar term is the standard Woods-Saxon potential

$$U_0(r) = -\frac{V}{1 + exp[(r - R)/a]}, \tag{13}$$

with the parameters $V = 52.03$ MeV, $R = 1.2709A^{1/3}$ fm, and $a = 0.742$ fm. Two other terms $U_{ls}(r)$ and $U_\tau(r)$ were obtained in a self-consistent way using the two-body Migdal-type interaction [17] for the spin-orbit and isovector part of the potential. The last term is the Coulomb potential of a uniformly charged sphere with $R_c = 1.18A^{1/3}$ fm.

In the next step, the residual strong interaction between the odd particle and the particles in the core should be taken into account. In IPM the state of an odd nucleus can be presented as

$$|\nu\rangle = a^\dagger_\nu |0\rangle,$$

where $|0\rangle$ is the ground state of an even core. When the interaction is switched on the odd nucleus state becomes a complicated superposition of the exited states of the core and the odd particle

$$|\nu) = Aa^\dagger_\nu|0\rangle + B_{\nu'\nu}(1p1h)a^\dagger_{\nu'}|1p1h\rangle + C_{\nu'\nu}(2p2h)a^\dagger_{\nu'}|2p2h\rangle + \cdots. \tag{14}$$

(Note that the round brackets mean a perturbed state). This superposition can be conveniently written as

$$|\nu) = \hat{U}a^\dagger_\nu|0\rangle, \tag{15}$$

where $\hat{U}$ is a unitary transformation depending on a nucleon-nucleon interaction. The matrix elements of any operator $\hat{O}$ between the exact states of the odd nucleus can be expressed by matrix elements of an effective operator $\hat{\tilde{O}}$ between the IPM states

$$(\nu'|\hat{O}|\nu) = \langle\nu'|\hat{\tilde{O}}|\nu\rangle, \tag{16}$$

where $\hat{\tilde{O}} = U^\dagger \hat{O} U$ is the effective operator. For a one body operator, like the Schiff moment, in first order in the nucleon-nucleon interaction the matrix elements of the effective operator coincide with Eq.(12) for the core contribution.

$$\langle\nu'|\tilde{O}|\nu\rangle = \langle\nu'|O|\nu\rangle + \sum_{\nu_1\nu_2} \langle\nu'\nu_1|V|\nu_2\nu\rangle \frac{n_{\nu_1} - n_{\nu_2}}{\epsilon_{\nu_1} - \epsilon_{\nu_2}} \langle\nu_2|O|\nu_1\rangle, \tag{17}$$

where V is the full nucleon-nucleon residual interaction. Leaving only 1p1h intermediate states we can sum over perturbation series obtaining for the effective Schiff moment the equation

$$\langle\nu'|\tilde{\mathbf{S}}|\nu\rangle = \langle\nu'|\mathbf{S}|\nu\rangle + \sum_{\nu_1\nu_2} \langle\nu'\nu_1|V|\nu_2\nu\rangle \frac{n_{\nu_1} - n_{\nu_2}}{\epsilon_{\nu_1} - \epsilon_{\nu_2}} \langle\nu_2|\tilde{\mathbf{S}}|\nu_1\rangle, \tag{18}$$

where $\langle\nu'|\mathbf{S}|\nu\rangle$ is the single particle matrix element of the bare Schiff moment operator, Eq.(1). Eq.(18) is just the RPA equation for the effective field in the terminology of Migdal's theory [17]. This equation describes the effects of core polarization due to interaction of the odd particle with the particles in the core.

Diagonal matrix elements of the Schiff moment are nonzero only in the presence of P and T violating interactions. For this reason the full nucleon nucleon interaction V must include this interaction

$$V(\mathbf{r}_1, \mathbf{r}_2) = F(\mathbf{r}_1, \mathbf{r}_2) + W(\mathbf{r}_1 - \mathbf{r}_2),$$

where F is the strong residual interaction and W is given by Eq.(5). In the same way, the mean field potential is the sum of two terms $U(r) + \delta U(\mathbf{r})$, where $U(r)$ is the main mean field created by the strong interaction and $\delta U(\mathbf{r})$ is the weak correction given by Eq.(10). As metioned above, $\delta U(\mathbf{r})$ creates the correction $\delta\psi_\nu(\mathbf{r})$ for every single particle wave function $\psi_\nu(\mathbf{r})$ in all matrix elements in Eq.(18). Since the corrections are small we can retain only the first order. To do this one should perform the substitutions $V \rightarrow F + W$ and all $\psi_\nu \rightarrow \psi_\nu + \delta\psi_\nu$ in Eq.(18) and keep the terms linear in W or $\delta\psi$. After this procedure the contributions to the Schiff moment can be written as a sum of three terms

$$S = \langle\delta\psi_\nu|\tilde{S}_z|\psi_\nu\rangle + \langle\psi_\nu|\tilde{S}_z|\delta\psi_\nu\rangle + \langle\nu|\delta S|\nu\rangle. \tag{19}$$

The first two terms are those where the weak correction enters via the odd particle wave function. $\tilde{S}$ satisfies the equation

$$\langle\nu'|\tilde{\mathbf{S}}|\nu\rangle = \langle\nu'|\mathbf{S}|\nu\rangle + \sum_{\nu_1\nu_2}\langle\nu'\nu_1|F|\nu_2\nu\rangle\frac{n_{\nu_1}-n_{\nu_2}}{\epsilon_{\nu_1}-\epsilon_{\nu_2}}\langle\nu_2|\tilde{\mathbf{S}}|\nu_1\rangle, \tag{20}$$

that differs from Eq.(18) only in the interaction. Here only the residual strong interaction F enters in the equation. The effects of the weak interaction are entirely in the wave functions of the odd proton. Eq.(20) describes the well known effect of renormalization of nuclear moments due to coupling with particle-hole states in the core. Renormalization of the Schiff moment due to coupling with the isoscalar dipole modes were discussed in [18]. It was found there that the effect of renormalization for ^{209}Bi nucleus is not significant, about 15%. We found similar values both for Cs and Fr nuclei for renormalization of the isoscalar component of the Schiff moment. The isovector componet, however, undergoes significant renormalization. Due to strong repulsion in the isovector channel it leads to reduction of the isovector component. As a result, the single particle contrubution becomes reduced by the factor ~ 2. For Cs, instead of values cited in the first line in Table I we obtain for the renormalized single particle contribution $\tilde{S}_{sp}$

$$\tilde{S}_{sp} = -0.003g_sg_0 - 0.056g_sg_1 - 0.016g_sg_2.$$

The third term in (19) satisfies the equation

$$\langle\nu'|\delta\mathbf{S}|\nu\rangle = \langle\nu'|\delta\mathbf{S}_0|\nu\rangle + \sum_{\nu_1\nu_2}\langle\nu'\nu_1|F|\nu_2\nu\rangle\frac{n_{\nu_1}-n_{\nu_2}}{\epsilon_{\nu_1}-\epsilon_{\nu_2}}\langle\nu_2|\delta\mathbf{S}|\nu_1\rangle, \tag{21}$$

that looks similar to (20). However, the inhomogenious term $\delta\mathbf{S}_0$ is completely different, namely

$$\begin{aligned}\langle\nu'|\delta\mathbf{S}_0|\nu\rangle = \sum_{\nu_1\nu_2}\langle\nu'\nu_1|W|\nu_2\nu\rangle\frac{n_{\nu_1}-n_{\nu_2}}{\epsilon_{\nu_1}-\epsilon_{\nu_2}}\langle\nu_2|\tilde{\mathbf{S}}|\nu_1\rangle+ \\ \sum_{\nu_1\nu_2}(\langle\nu'\delta\psi_{\nu_1}|F|\nu_2\nu\rangle + \langle\nu'\nu_1|F|\delta\psi_{\nu_2}\nu\rangle)\frac{n_{\nu_1}-n_{\nu_2}}{\epsilon_{\nu_1}-\epsilon_{\nu_2}}\langle\nu_2|\tilde{\mathbf{S}}|\nu_1\rangle+ \\ \sum_{\nu_1\nu_2}\langle\nu'\nu_1|F|\nu_2\nu\rangle\frac{n_{\nu_1}-n_{\nu_2}}{\epsilon_{\nu_1}-\epsilon_{\nu_2}}(\langle\delta\psi_{\nu_2}|\tilde{\mathbf{S}}|\nu_1\rangle + \langle\nu_2|\tilde{\mathbf{S}}|\delta\psi_{\nu_1}\rangle).\end{aligned} \tag{22}$$

The first term on the rhs of Eq.(22) is the core contribution (12), where instead of the bare Schiff moment operator (1) enters the renormalized operator $\tilde{\mathbf{S}}$. The second and the third terms correspond to additional contributions where the weak interaction enters via corrections to the intermediate single particle states $|\nu_1>$ and $|\nu_2>$. Equations (20-22)

TABLE IV: Effects of core polarization on the Schiff moment for odd proton nuclei, for finite range weak interaction. The bare values of the Schiff moment, without core polarization, are listed in the first line for each nucleus. The units are $e\,fm^3$

	$g_s g_0$	$g_s g_1$	$g_s g_2$
^{133}Cs	0.08	-0.02	-0.21
^{133}Cs	0.006	-0.02	-0.04
^{223}Fr	-0.122	-0.052	0.300
^{223}Fr	-0.009	-0.016	0.030

describe all of the core polarization effects. The main contribution to $\delta\mathbf{S}_0$ comes from the first term in Eq. (22). This term depends on the renormalized Schiff moment $\tilde{\mathbf{S}}$, therefore its contribution becomes reduced proportionally to the reduction of $\tilde{\mathbf{S}}$, approximately by the factor ~ 2. The contribution of the second and the third terms in Eq.(22) is small. It is noticable only in T=1 channel.

In Tables IV, and V we present the results of calculations of the Schiff moment for sets of odd proton and odd neutron nuclei. The effects of the core polarization are twofold. First, without core polarization, for odd neutron nuclei the Schiff moment depended only on one constant η_{np} (see Eq.(7)). Now, the Schiff moment depends on all three constants. Second, the values of the Schiff moment are decreased. There are two sources of that reduction. First, it is mentioned above reduction in the first term in Eq.(22). Further decrease comes from repulsive nature of spin-spin residual forces responsible for renormalization of $\delta\mathbf{S}_0$ in Eq.(21).

IV. CONCLUSIONS

We stress here two main conclusions from this work. First, the core-polarization contributions to the Schiff moments are sizeable. These corrections reduce the lowest order values obtained in the past by factors two to three. This has some important consequences when trying to interpret the implication of the experimental atomic dipole measurements. The reductions of the Schiff moments (compared to the values obtained in the past) means

TABLE V: Effects of core polarization on the Schiff moment for neutron odd nuclei, for finte range weak interaction. The bare values of the Schiff moment, without core polarization, are listed in the first line for each nucleus. The units are $e\,\mathrm{fm}^3$.

	$g_s g_0$	$g_s g_1$	$g_s g_2$
^{199}Hg	-0.09	-0.09	0.18
^{199}Hg	-0.00004	-0.055	0.009
^{129}Xe	0.06	0.06	-0.12
^{129}Xe	0.008	0.006	-0.009
^{211}Rn	-0.12	-0.12	0.24
^{211}Rn	-0.019	0.061	0.053
^{213}Ra	-0.12	-0.12	0.24
^{213}Ra	-0.012	-0.021	0.016
^{225}Ra	0.08	0.08	-0.16
^{225}Ra	0.033	-0.037	-0.046

that using the experimental upper limits for the atomic dipole one obtains now a higher upper limit for the strength parameters in the weak parity and time reversal non-conserving component of the N-N interaction. Second, we have seen that the finite range reduces the mixing matrix elements between the opposite parity states in the case of spherical nuclei. It was argued [7] that the large enhancement factors for the Schiff moments found in quadrupole+octupole nuclei [11] will be reduced when finite range interactions are used. We see that the enhancement factors (compared to the spherical case) will not be reduced because the effect of finite range has a similar effect in spherical and in deformed nuclei.

We wish to thank V. Zelevinsky for helpful discussions. One of us (V.D.) wants to thank Tel-Aviv University for the hospitality. This work was supported by the US – Israel Binational Science Foundation.

---

[1] J.P. Jacobs, W.M. Klipstein, S.K. Lamoreaux, B.R. Heckel, E.N. Fortson Phys. Rev. A**52**, 3521(1995)

[2] V.V. Flambaum, I.B. Khriplovich and O.P. Sushkov, Nucl.Phys. **A**449, 750, (1986)]

[3] M.V. Romalis, W.C. Griffith, J.P. Jacobs, E.N. Fortson, Phys. Rev. Lett. **86**, 2505 (2001).

[4] L.I. Schiff, Phys. Rev. **132**, 2194 (1963).

[5] V.V. Flambaum, I.B. Khriplovich, and O.P. Sushkov Zh. Eksp. Teor. Fiz. **87**, 1521 (1984) [Sov. Phys. JETP **60**, 873 (1984)].

[6] V.F. Dmitriev and R.A. Sen'kov, Yad.Fiz. **66** 1988 (2003), (Phys.Atom.Nucl. **66**, 1940 (2003)), nucl-th/0304048.

[7] J. Engel, M. Bender, J. Dobaczewski, J.H. de Jesus, and P. Olbratowski, Phys.Rev. C**68**, 025501 (2003).

[8] P. Möller, J. R. Nix, W. D. Myers, and W. J. Swiatecki, Atomic Data Nucl. Data Tables **59**, 185 (1995).

[9] P. Möller, J. R. Nix, and K.-L. Kratz, Atomic Data Nucl. Data Tables **66**, 31 (1997) .

[10] N. Auerbach, V.V. Flambaum, V. Spevak, Phys.Rev.Lett. **76**, 4316 (1996).

[11] V. Spevak, N. Auerbach, V.V. Flambaum, Phys. Rev. C **56**, 1357, (1997).

[12] G. Barton, Nuovo Cim., **19**, 512 (1961).

[13] W. C. Haxton and E. M. Henley, Phys. Rev. Lett., **51**, 1937 (1983).

[14] 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**, 77 (1988).

[15] I.B. Khriplovich, *Parity Nonconservation in Atomic Phenomena*, (Gordon and Breach, Philadelphia, 1991)

[16] O.P. Sushkov, V.V. Flambaum and I.B. Khriplovich, Zh. Eksp. Teor. Fiz. **87**, 1521 (1984); [Sov. Phys. JETP, **60** 873 (1984)].

[17] A.B. Migdal, Theory of Finite Fermi Systems and Applications to Atomic Nuclei, Wiley, New York, 1967.

[18] I. Hamamoto and B.A. Brown, Phys. Rev.C **62**, 024318, (2000).

[19] B.L. Birbrair and V.A. Sadovnikova, Yad. Fiz. **20**, 645 (1974) [Sov. J. Nucl. Phys. **20**,347 (1974).