P- and T-violating Schiff moment of the Mercury nucleus

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Abstract

The Schiff moment of the $^{199}$Hg nucleus was calculated using finite range P- and T-violating weak nucleon-nucleon interaction. Effects of the core polarization were considered in the framework of RPA with effective residual forces.

1 Introduction

The most precise limit on parity- and time-invariance violating nucleon-nucleon interaction has been obtained from the measurement of the atomic electric dipole moment of $^{199}$Hg [1]. The hadronic part of the atomic dipole moment associated with the electric dipole moment of the $^{199}$Hg nucleus manifests itself through the Schiff moment which is the first nonzero term in the expansion of the nuclear electromagnetic potential after including the screening of the atomic electrons [2, 3, 4].

The operator for the Schiff moment is [5]

$$S_\mu = \frac{1}{10} \sqrt{\frac{4\pi}{3}} \sum_i e_i \left( \begin{array}{c} r_i^3 - \frac{5}{3} \langle r_i^2 \rangle \end{array} \right) Y_{1\mu}(\hat{r}_i),$$

where $e_i$ is $e$ for a proton and zero for a neutron. The Schiff moment generates T- and P-odd electrostatic potential in the form

$$\phi(r) = 4\pi S \cdot \nabla \delta(r).$$

Interaction of atomic electrons with the potential given by Eq. 2 produces an atomic dipole moment

$$d_{\text{atom}} = \sum_n \frac{|0\rangle - e \sum_i Z_i \phi(r_i)|n\rangle - e \sum_i Z_i |0\rangle}{E_n - E_0} + \text{h.c.}$$

Due to contact origin of the potential, the electrons in $s$- and $p$- atomic orbitals only contribute to the dipole moment given by Eq. 3.
The nuclear Schiff moment was calculated so far in a simplified model [5,6] without considering many-body nuclear structure effects. These effects have to be understood properly if we intend to extract the parameters of P- and T-violating nuclear interaction from the value of the Schiff moment. For light nuclei the properties of the Schiff moment strength obtained in modern shell model calculations were discussed in Ref. [7]. The study of the polarization effects associated with the coupling to isoscalar dipole compression mode has been performed in Ref. [8]. In our paper we calculate the Schiff moment of $^{199}$Hg nucleus within RPA framework with effective residual strong forces using finite range P- and T-odd weak nuclear interaction.

2 Basic ingredients of the theory

2.1 Nucleon-nucleon P- and T-odd interaction

We use the interaction generated by P- and T-violating pion exchange [9,10,11].

$$W(r_1 - r_2) = -\frac{g}{8\pi m_p} \left[ (g_0 \tau_1 \cdot \tau_2 + g_2 (\tau_1 \cdot \tau_2 - 3\tau_1^3 \tau_2^3)) (\sigma_1 - \sigma_2) \right. $$

$$\left. + g_1 (\tau_1^3 \tau_3^3 - \tau_2^3 \tau_3^3) \cdot \nabla_1 e^{-m_\pi r_{12}} \right] \cdot \frac{r_{12}}{r_{12}},$$

(4)

where $g$ is the usual strong pion-nucleon pseudoscalar coupling constant, $g_0$, $g_1$, and $g_2$ correspond to isoscalar, isovector, and isotensor P- and T-odd couplings, $m_p$ is the proton mass. In contrast to P-odd and T-even interaction, in Eq. (4) the exchange of $\pi^0$ is allowed. This term produces the direct contribution to P- and T-odd part of nuclear mean field while the other terms produce the exchange contribution only. Since the direct contribution dominates for finite range potentials we can expect that the interaction (4) is the leading one and the exchange of heavier mesons can be omitted.

In previous calculations the phenomenological contact interaction has often been used instead of finite range interaction given by Eq. (4). It has the form [12]. It has the form

$$W_c(r_a - r_b) = \frac{G}{\sqrt{2} m_p} \left( (\eta_{ab} \sigma_a - \eta_{ab} \sigma_b) \cdot \nabla_a \delta(r_a - r_b) + \eta_{ab} (\sigma_a \times \sigma_b) \cdot \{(p_a - p_b) \cdot \delta(r_a - r_b)\} \right),$$

(5)

where $G$ is the Fermi constant. In the limit $m_x \to \infty$ the interaction transforms into Eq. (4) after the substitution $g g_i \rightarrow \frac{G m^2}{\sqrt{2}} \eta$. We shall use this factor when comparing our results with those obtained using the contact interaction given by Eq. (5).

2.2 Nuclear mean field and correction from the weak forces

In our calculations we used full single-particle spectrum including continuum. The single-particle basis was obtained using partially self-consistent mean-field potential of Ref. [13]. The potential includes four terms. The isoscalar term is the standard Woods-Saxon potential

$$U_0(r) = -\frac{V}{1 + \exp \frac{r - R}{a}},$$

(6)
with the parameters $V = 52.03$ MeV, $R = 1.2709A^{1/3}$ fm, and $a = 0.742$ fm.

Two other terms $U_{is}(r)$, and $U_{e}(r)$ were calculated self-consistently using two-body Migdal-type interaction of Ref. [14] for the spin-orbit and isovector parts of the potential. The last term is the Coulomb potential calculated for uniformly charged sphere with $R_C = 1.18A^{1/3}$ fm. The mean field potential obtained in this way produces good fit for single particle energies and r.m.s. radii for nuclei in the lead region.

The correction to the mean field (6) from the weak interaction (4) consists of direct and exchange terms. The direct term has the form

$$\delta U_{dir}(r) = \frac{gM_n^2}{\pi m_p} (\sigma \cdot n) r^3 \int_0^\infty r'^2 dr' b_{10}(r,r') [(g_0 - 2g_2)(\rho_p(r) - \rho_n(r)) + g_1(\rho_p(r) + \rho_n(r))],$$

(7)

where the function $b_{10}(r_1, r_2)$ is a combination of spherical Bessel functions of imaginary argument

$$b_{1l2}(r_1, r_2) = i l_1 (m_\pi r_1) k_l (m_\pi r_2) \theta(r_2 - r_1) - i l_2 (m_\pi r_2) k_l (m_\pi r_1) \theta(r_1 - r_2).$$

(8)

Note, that the potential given by Eq. (7) is pure isovector. The contribution of the isovector interaction component dominates in Eq. (7), the isoscalar and isotensor components of the interaction are suppressed by the factor $N/A$. For zero range interaction the potential would be proportional to $\nabla \rho(r)$. The gradient makes this potential very sensitive to the details of nuclear surface. Our potential given by Eq. (4) is less sensitive to the surface due to additional integration over region of the order of pion Compton wavelength.

The exchange term is more complicated. The matrix element of it takes the form

$$\langle \hat{\varphi} | \delta U_{exch}(r,r') | \nu \rangle = W_{\nu}(r,r'),$$

where $|\hat{\varphi}\rangle = -(\sigma \cdot n)|\nu\rangle$, and

$$W_{\nu}(r,r') = \frac{1}{2\nu + 1} \frac{gM_n^2}{\pi m_p} TR_{2} \left\{ \sum_{\kappa l_1 l_2} \left( \frac{g_0}{2} (3 - r_1^3 r_2^3) - 2g_2 r_1^3 r_2^3 + \frac{g_1}{2} (r_1^3 + r_2^3) \right) \times \begin{pmatrix} l_1 \\
0 \\
0 \\
0 \end{pmatrix} b_{1l2}(r,r') R_{\kappa}(r) R_{\kappa}(r') \left[ (-)^{l_1} [l_1] (|\kappa| T_{l_2}^{l_1} |\nu\rangle + |\nu| T_{l_2}^{l_1} |\kappa\rangle) - (-)^{l_2} [l_2] (|\kappa| T_{l_2}^{l_1} |\nu\rangle + |\nu| T_{l_2}^{l_1} |\kappa\rangle) \right] \right\}.$$

(9)

The Trace is assumed over isospin variable of the second particle, $n_\kappa$ is the occupation number of the single particle state $|\kappa\rangle$, $[l] = \sqrt{2l+1}$, and the tensor operator $T_{JM}^{l} (n) = \{ \sigma \otimes Y_L(n) \}_{J M}$.

The correction to the single particle wave function $\psi_\nu(r)$ can be presented as

$$\delta \psi_\nu(r) = (\sigma \cdot n) \Omega_\nu(n) \delta R_{\nu}(r),$$

where $\Omega_\nu(n)$ is the angular part of the wave function. The radial correction $\delta R_\nu(r)$ is the sum of the direct and the exchange terms $\delta R_\nu(r) = \delta R_{\nu \ dir}(r) + \delta R_{\nu \ exch}(r)$, where

$$\delta R_{\nu \ dir}(r) = \int_0^\infty G_{\nu l_1}(r, r' |\epsilon_\nu) \delta U_{dir}(r') R_{\nu}(r') r'^2 dr',$$
The operator interaction only is included. The corrections from the weak forces for these terms arising from the P- and T-violating corrections to the wave functions of an odd nucleon only. The third term presents an element of the Schiff moment as a sum of three terms the terms linear in $\delta \psi$. For static moments the interaction in the vacuum. We do not discuss this effect here and keep the weak interaction in the form (4). Strictly speaking, the interaction of two nucleons in nuclear matter differs from the interaction in the vacuum. Here $G_{\mu\nu}(r, r'|\epsilon_\nu)$ is the Green function of radial Schrödinger equation for the total angular momentum $j_\nu$, and the orbital angular momentum $I_\nu = 2j_\nu - l_\nu$. $\epsilon_\nu$ is the single particle energy.

### 3 Core polarization

The effects of the core polarization for a one particle operator can be treated introducing a renormalized operator $\tilde{S}$ satisfying the equation

$$\tilde{S}_{\nu'\nu} = S^0_{\nu'\nu} + \sum_{\mu'\mu} \tilde{S}_{\mu\mu'} \frac{n_\mu - n_{\mu'}}{\epsilon_\mu - \epsilon_{\mu'}} \langle \nu'\mu'|F + W|\mu\nu\rangle,$$

where $S^0$ is the bare Schiff moment operator given by Eq. (11), $n_\mu$ and $\epsilon_\mu$ are the single particle occupation numbers and energies. For static moments the external frequency $\omega \to 0$. The interaction in Eq. (11) includes both the strong residual interaction $F$ and the weak one. The latter we take in the form given by Eq. (4). Strictly speaking, the interaction of two nucleons in nuclear matter differs from the interaction in the vacuum. We do not discuss this effect here and keep the weak interaction in the form (4). The single particle wave functions in Eq. (4) are the eigenstates of the mean field which is also the sum of the strong and the weak fields given by Eqs. (7,9). Since the weak forces are really small compared to the strong interaction, it is natural to treat them perturbatively. The simplest way to do it is to present $\langle r|\nu \rangle = \psi_\nu(r) + \delta \psi_\nu(r)$ and to gather the terms in $\delta \psi$ and $W(r_1 - r_2)$. It is convenient to present the matrix element of the Schiff moment as a sum of three terms

$$\tilde{S}_{\nu'\nu} = \langle \delta \psi_\nu|S|\psi_\nu\rangle + \langle \psi_\nu|S|\delta \psi_\nu\rangle + \langle \psi_\nu|\delta S|\psi_\nu\rangle.$$

The operator $S$ in the first two terms satisfies the same Eq. (4) where the strong interaction only is included. The corrections from the weak forces for these terms are in the wave functions of an odd nucleon only. The third term presents an induced contribution arising from the P- and T-violating corrections to the intermediate states $|\mu\rangle$ and $|\mu'\rangle$. The equation for $\delta S$ is

$$(\delta S - \delta S_{NL})_{\nu'\nu} = (\delta S_0)_{\nu'\nu} + \sum_{\mu'\mu} (\delta S - \delta S_{NL})_{\mu\mu'} \frac{n_\mu - n_{\mu'}}{\epsilon_\mu - \epsilon_{\mu'}} \langle \nu'\mu'|F|\mu\nu\rangle,$$

where $\delta S_{NL}$ is the non-local part of $\delta S$ produced by the exchange matrix elements of weak interaction $W$

$$(\delta S_{NL})_{\nu'\nu} = \sum_{\mu'\mu} (S)_{\mu\mu'} \frac{n_\mu - n_{\mu'}}{\epsilon_\mu - \epsilon_{\mu'}} \langle \nu'\mu'|W_{exch}|\mu\nu\rangle.$$

The equation for $\delta S_0$ is

$$(\delta S_0)_{\nu'\nu} = \sum_{\mu'\mu} \frac{n_\mu - n_{\mu'}}{\epsilon_\mu - \epsilon_{\mu'}} \langle \mu|S|\mu'\rangle \langle \nu'\mu'|W_{dir}|\mu\nu\rangle.$$

Here $G_{\mu\nu}(r, r'|\epsilon_\nu)$ is the Green function of radial Schrödinger equation for the total angular momentum $j_\nu$, and the orbital angular momentum $I_\nu = 2j_\nu - l_\nu$. $\epsilon_\nu$ is the single particle energy.
\begin{equation}
\langle \mu | S | \mu' \rangle \langle \nu' | \delta \psi_{\nu'} | F | \mu \nu \rangle + \langle \mu | S | \mu' \rangle \langle \nu' | \mu' | F | \delta \psi_{\nu'} | \rangle.
\end{equation}

Note, that in the absence of the core polarization the first term only in Eq. \ref{eq:16} contributes into $\delta S_0$. And this is the only term that produces the Schiff moment of the nucleus with an odd neutron, like $^{199}$Hg. The residual interaction $F$ has the form

\begin{equation}
F = C (f(r) + f'(\tau \cdot \tau_2) + g_s(\sigma_1 \cdot \sigma_2) + g_s'(\sigma_1 \cdot \sigma_2)(\tau_1 \cdot \tau_2)) \delta(r_1 - r_2),
\end{equation}

where $C = 300 \text{ MeV fm}^3$, and $f(r) = f_{\text{ex}} + (f_{\text{in}} - f_{\text{ex}}) \frac{\rho(r)}{\rho(0)}$. Note that the angular dependence of the operators $S$ and $\delta S$ is completely different. While $S_{\mu} \sim Y_{1 \mu}(n)$, the induced part $\delta S_{\mu}$ is a superposition of spin dependent operators $\sigma_\mu$ and $\{\sigma \otimes Y_2(n)\}_1\mu$. For this reason, different parts of the interaction \ref{eq:16} contribute into renormalization of $S$ and $\delta S$.

For a spherical nucleus we can separate the angular variables and solve the obtained equations in coordinate space. The equations are

\begin{equation}
S^a(r) = S^a_0(r) + \int_0^{\infty} A^{ab}(r, r') S^b(r') \, dr',
\end{equation}

where $a = p, n$ and $S^a_0(r)$ is the radial part of the Schiff moment operator Eq. \ref{eq:11} multiplied by $r$. The particle-hole propagator $A(r, r')$ was calculated by means of the Green functions of radial Schrödinger equation.

\begin{equation}
A(r, r') = \frac{C}{3} \cdot T_{2l}\{(f(r) + f'(\tau \cdot \tau_2)) \sum_{\kappa j l} n_{\kappa l} |j l||\kappa\rangle^2 r R_\kappa(r) \times \nonumber \end{equation}

\begin{equation}
r' R_\kappa(r') \{G_{jl}(r, r'|\epsilon_\kappa + \omega) + G_{jl}(r, r'|\epsilon_\kappa - \omega)\}.
\end{equation}

Similar equations can be written for $\delta S(r)$ - the local part of the induced moment. They differ from Eq. \ref{eq:16} in type of tensor operators and in residual interaction.

### 3.1 Separation of the spurious component

The integral equation \ref{eq:17} must have zero eigenmode related to the center of mass motion. However, in our case the mean field Eq. \ref{eq:10} is not consistent with the interaction Eq. \ref{eq:16} and, in general, we do not have zero energy for the center of mass motion. The situation can be improved using some freedom in the value of the interaction constant $f_{\text{in}}$ in Eq. \ref{eq:10}. We can fix the value by the condition $\omega_0 = 0$ for the lowest isoscalar dipole mode. Since this procedure is numerical the condition $\omega_0 = 0$ cannot be fulfilled exactly but with finite accuracy. For $^{199}$Hg we found that the set $f_{\text{in}} = 0.3935, f_{\text{ex}} = -2.6$, and $f' = 1.07$ gives for $\omega_0$ the value $\omega_0 = 0.1 \text{ KeV}$ which is really small compared to the energy of dipole transitions.

The finite accuracy in determination of the spurious mode brings another problem for solutions of the Eq. \ref{eq:17}. The bare operator $S_0(r)$ becomes non orthogonal to the spurious mode transition density. This results in admixture of the spurious component to all solutions of the Eq. \ref{eq:14}. The spurious component should be subtracted since it can change the solution considerably. The subtraction can be performed using analytical properties of the solution as a function of $\omega$. The solution of the inhomogeneous linear integral equation \ref{eq:14}
as a function of $\omega$ has the first order poles at $\omega = \pm E_{ex}$, where $E_{ex}$ are the excitation energies of the RPA modes. Thus, at small $\omega$ the solution can be presented as

$$S(r|\omega) = \frac{a(r)}{\omega^2 - \omega_0^2} + b(r|\omega = 0) + O\left(\frac{\omega^2}{E_{ex}}\right).$$

(19)

The first term in Eq. (19) is the spurious component contribution. Let us define the following set of integrals in a complex $\omega$-plane along a circle with the radius satisfying the conditions $\omega_0 \ll |\omega| \ll E_{ex}$.

$$J_n(r) = \oint \omega^n S(r|\omega) \frac{d\omega}{2\pi i}.$$  

(20)

Using Eq. (19) we can perform the integration analytically. As a result we obtain

$$b(r|\omega = 0) = J_{-1}(r),$$

$$\omega_0^2 = J_3(r)/J_1(r).$$

(21)

The numerical integration in Eq. (20) has been performed using 32-points Gauss formula. The integration radius was $|\omega| = 0.1$ MeV. The obtained field $b(r|\omega = 0)$ was practically insensitive to the integration radius. The value of $\omega_0^2$ was independent on $r$ in first 5 or 6 digits. The loss of accuracy was noticeable only at $|\omega| \sim \omega_0$.

Fig.1 shows the result for renormalized $S(r) = b(r|\omega = 0)$ with subtracted spurious component. It is shown by full line. The dashed line shows the unrenormalized component $S_0(r)$. The effects of the core polarization are not large. We found that for $^{209}$Bi they change the valence proton contribution by $\sim 15\%$ which is in fair agreement with the estimates of Ref. [8]. The proton component does not contribute to the Schiff moment of $^{199}$Hg since the valence nucleon is a neutron. However, due to neutron-proton residual strong interaction some neutron component is produced by the core polarization. This component is shown in Fig.2. Qualitatively, the behavior of the neutron component inside a nucleus resembles the behavior of the bare operator in Fig.1. However, the magnitude of the neutron component is smaller and the behavior outside of the nucleus is completely different.

4 Results for $^{199}$Hg

In Table 1 we show the contributions of the neutron component discussed above to the Schiff moment of $^{199}$Hg. The three columns correspond to three isospin channels. The first row is the contribution from $\delta R_{dir}$ produced by $\delta U_{dir}$ given by Eq. (7). The second row is the contribution from $\delta R_{exch}$ produced by $\delta U_{exch}$ given by Eq. (9), and the third row is the sum of them.

In the direct contributions the isospin channel $T = 1$ dominates, as it was mentioned above. The exchange contributions are small, in general. For the isospin channels $T = 0, 2$ they are comparable to the direct contributions just because the latter’s are suppressed by the factor $(N - Z)/A$.

The contributions from $\delta S$ are more significant. In Table 2 we list the contributions from the direct and the exchange parts of the weak interaction $W$. In the first two rows the unrenormalized contributions $\delta S_0$ are shown. Again,
Figure 1: The proton component of the Schiff moment. Full line is the renormalized operator after subtraction of the spurious component. Dashed line is the bare operator Eq. (1).

Figure 2: The neutron component of the renormalized Schiff moment.

Table 1: Contributions of the neutron component to the Schiff moment of $^{199}$Hg in $[\text{e fm}^3]$.

|       | $gg_0$ | $gg_1$ | $gg_2$ |
|-------|--------|--------|--------|
| direct| 0.0038 | -0.024 | -0.0076|
| exchange| 0.0032 | -0.002 | -0.0015|
| total  | 0.0070 | -0.026 | -0.0091|

the exchange contributions are considerably smaller than the direct ones. The contributions from the channels $T = 1$, and $T = 2$ are comparable here although $T = 1$ contribution is still larger. Next two rows show the renormalized contributions $\delta S$. The effect of the core polarization is significant here. The induced moment $\delta S$ includes spin dependent operators, therefore, the spin-spin part of the residual interaction Eq. (16) is responsible for their renormalization. The spin-spin interaction is repulsive with the constants $g_s = 0.63$, and $g_s' = 1.01$. The repulsion results in decrease of absolute values of the renormalized contributions. Finally, in the last row we show the contribution of the non-local term $\delta S_{NL}$ Eq. (15). It has the exchange origin as well and its contribution is really insignificant. All the contributions can be summarized as follows

$$ S = -0.0004gg_0 - 0.055gg_1 + 0.009gg_2 \ [\text{e fm}^3]. $$

(22)

The obtained value for the Schiff moment in Eq. (22) cannot be compared directly with the previous calculations with the contact interaction Eq. (5). The reason is in different definition of the dimensionless constants $gg_i$ in Eq. (4) and $\eta_{ab}$ in Eq. (5). To perform the comparison we redefine the constants $g_i$

$$ g_i = \frac{Gm^2_{\pi}}{\sqrt{2}} \tilde{g}_i. $$

(23)
Table 2: Induced contributions to the Schiff moment of $^{199}$Hg in [e fm$^3$].

|         | $g\bar{g}_0$ | $g\bar{g}_1$ | $g\bar{g}_2$ |
|---------|--------------|--------------|--------------|
| direct $\delta S_0$ | -0.0302      | -0.0631      | 0.0604       |
| exchange $\delta S_0$ | -0.0007      | -0.0012      | -0.0007      |
| direct $\delta S$ | -0.0086      | -0.0285      | 0.0172       |
| exchange $\delta S$ | -0.0002      | -0.0008      | -0.0003      |
| non-local $\delta S_{NL}$ | 0.0014       | -0.00004     | 0.0013       |

With this factor the integration over space of the Yukawa function gives 1, exactly as the integration of a $\delta(r)$. Introducing this factor we obtain

$$S = (-0.01g\tilde{g}_0 - 0.86g\tilde{g}_1 + 0.14g\tilde{g}_2) \times 10^{-8} \text{ [e fm}^3\text{]}.$$

(24)

This value should be compared with $S = -1.4 \times 10^{-8} \eta_{np}$ from Ref. [10], and $S \approx -1.6 \times 10^{-8} \eta_{np}$ from Ref. [6]. Remembering that $\eta_{np} \sim g(\tilde{g}_0 + \tilde{g}_1 - 2\tilde{g}_2)$ we conclude that the difference between our result and previous calculations is significant for $T = 0$, and $T = 2$ channels. Our values are smaller in absolute value. In order to trace the origin of this difference we repeated our calculations using the contact interaction and omitting the core polarization. The contact interaction was obtained by replacing the Yukawa function in Eq. (4) by the delta-function. The result is

$$S = -0.96 \times 10^{-8} g(\tilde{g}_0 + \tilde{g}_1 - 2\tilde{g}_2) \text{ [e fm}^3\text{]}.$$

(25)

If we omit completely the effect of the core polarization in calculations with the finite range interaction Eq. (4) then the only nonzero contribution in the first row of the Table 2 becomes

$$S = -0.086 g(\tilde{g}_0 + \tilde{g}_1 - 2\tilde{g}_2) \text{ [e fm}^3\text{]},$$

that corresponds to

$$S = -1.35 \times 10^{-8} g(\tilde{g}_0 + \tilde{g}_1 - 2\tilde{g}_2) \text{ [e fm}^3\text{]}.$$

(26)

Comparing Eq. (25) and Eq. (26) we conclude that the effect of finite weak interaction range is not very significant. The main effect bringing the value of the Schiff moment from that in Eq. (26) to the value in Eq. (24) comes from the core polarization.

The last remark concerns the pairing effects. The nucleus $^{199}$Hg has 7 neutron holes in the unfilled shell. One of them is fixed in $p_{1/2}$ state and 6 others should be distributed among the other states in the shell. From mass differences we found $\Delta_n = 0.69$ MeV. This value of the pairing gap is typical for developed pairing. For dipole transitions the transition energy is large compared to $\Delta$ and the pairing effects are small. They were omitted in Eq. (17). For the induced moment the situation is different. There, the transitions with $\Delta J = 0, 1$ and $\Delta L = 0, 2$ are responsible for the core polarization. Such transitions exist inside the last unfilled neutron shell and, due to Pauli blocking, they are sensitive to
the details of the shell occupation. For a T-odd operator the effects of pairing in
the core polarization can be considered in the way used in Ref. [15]. We found,
that the main effect of pairing is in fixing the occupation numbers in the upper
unfilled neutron shell. As soon as we keep the occupation numbers fixed the
results for $\delta S$ are changed within few percents when we put $\Delta_n = 0$.

In summary, we calculated the Schiff moment of $^{199}\text{Hg}$ nucleus using finite
range weak interaction and considering the core polarization effects. The effects
of the finite interaction range are not very significant for the Mercury nucleus.
They do not change the order of magnitude of the Schiff moment calculated
with the contact interaction. The effects of the core polarization are two-fold.
First, they renormalize the bare operator of the Schiff moment producing a small
neutron component due to n-p residual interaction. Second, they produce an
induced Schiff moment due to P- and T- violation component in the intermediate
single particle states. The induced moment is proportional to both the strong
residual and the weak interactions. The effects of the core polarization for the
Mercury nucleus are large and they have to be accounted in calculations of P-
and T-violating effects in the Mercury nucleus.

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