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ИНСТИТУТ ЯДЕРНОЙ ФИЗИКИ СО АН СССР

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RELATIVISTIC MANY-BODY CALCULATION OF PARITY NON-CONSERVATION IN THALLIUM

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ABSTRACT

Parity non-conserving (PNC) E1-amplitudes for ^{205}Tl are calculated. In units of $10^{-10}\cdot(i|e|a_B Q_W/-N)$ (Q_W is the weak charge of the nucleus, N is the number of the neutrons)

$$\langle 6p_{3/2}\frac{1}{2}|D_z|6p_{1/2}\frac{1}{2}\rangle = -2.70\cdot(1\pm0.03),$$

$$\langle 7p_{1/2}\frac{1}{2}|D_z|6p_{1/2}\frac{1}{2}\rangle = -0.81\cdot(1\pm0.06).$$

Using the developed technique we recalculate the PNC E1-amplitude for ^{133}Cs with high accuracy:

$$\langle 7s\frac{1}{2}|D_z|6s\frac{1}{2}\rangle = 0.089\cdot(1\pm0.02).$$

We calculate also the PNC E1-amplitude for ^{85}Rb:

$$\langle 6s\frac{1}{2}|D_z|5s\frac{1}{2}\rangle = 0.0138\cdot(1\pm0.02).$$

For calculation we use time-dependent Hartree—Fock method as the first approximation and then take into account all the correlation corrections of second order in residual Coulomb interaction.

1. INTRODUCTION

For the thallium atom parity non-conservation (PNC) in forbidden M1-transition $6p_{1/2}\to 7p_{1/2}$ has been measured in Refs [1—5]. There are many calculations of PNC E1-amplitude for this transition [6—12]. In Refs [6, 7] the calculation was done by semiempirical method. The authors of Ref. [8] used the method of effective potential. The relativistic many-body perturbation theory was used in Refs [9—12]. The PNC in transition $6p_{1/2}\to 6p_{3/2}$ is not investigated experimentally yet. The previous calculations were carried out in Refs [13—15, 8]. This transition is similar to the M1-transitions in Bi for which the PNC has been measured. It is well known that for Bi there is essential disagreement between results of different calculations of PNC effect. We think that high presision calculation of PNC effect for $6p_{1/2}\to 6p_{3/2}$ transition in thallium is a good experience for Bi.

In the present work the calculations of PNC E1-amplitudes are carried out by the methods of relativistic many-body perturbation theory:

1. As zero approximation we use the wave-functions of $V^{(N-1)}$ relativistic Hartree-Fock (RHF) method.

2. At the second step we take into account the polarization of closed atomic shells by PNC weak interaction and by the external electromagnetic wave. For this we use time-dependent Hartree-Fock (TDHF) method which is equivalent to the random phase approximation (RPA).

3. The final and most complicated step is the calculation of all the correlation corrections of the second order in residual Coulomb interaction.

This way of calculation is similar to way used by us for calculation of PNC $6s \to 7s$ E1-amplitude for Cs [16]. However the technique of the present work has the important distinctions from that of Ref. [16]. First, the electromagnetic polarization is taken into account here in all the orders in residual Coulomb interaction (in Ref. [16] in the first order only). Second and most important distinction is the calculation of all the second order correlation corrections. In Ref. [16] only the dominating correlation diagrams which correspond to renormalization of wave-functions, were taken into account. Besides the calculation for ^{205}Tl we use the developed technique to calculate all the small second order contributions to PNC E1-amplitude for ^{133}Cs which were not taken into account in Ref. [16]. This allows us to decrease the theoretical uncertainty of calculation. We calculate as well the PNC E1-amplitude for $5s \to 6s$ transition in ^{58}Rb.

In Refs [11, 12] the PNC E1$(6p_{1/2} \to 7p_{1/2})$-amplitude for ^{205}Tl has been calculated by the TDHF-method. Our corresponding value (the points 1 and 2 of the above program) is in agreement with the results of Refs [11, 12].

2. TIME-DEPENDENT HARTREE-FOCK CALCULATION (TDHF)

Let us remind the reader the scheme of $V^{(N-1)}$-approximation for the atoms with one unpaired electron. To find the wave-functions one should solve in self-consistent way the Hartree-Fock equations for the closed shells

$$(H_0 - \varepsilon_k)\psi_k = 0, \tag{1}$$

$$H_0 = \vec{\alpha}\vec{p} + (\beta - 1)m - \frac{Ze^2}{r} + V^{(N-1)}, \tag{2}$$

$$V^{(N-1)}\psi = \sum_{i=1}^{N-1} \left\{ \int \psi_i^+(r_1)\,\psi_i(r_1)\,\frac{e^2}{r_{12}}\psi(r_2)\,d^3r_1 - \right.$$

$$\left. - \int \psi_i^+(r_1)\,\psi(r_1)\,\frac{e^2}{r_{12}}\psi_i(r_2)\,d^3r_1 \right\}. \tag{3}$$

Here $\vec{\alpha}$ and β are Dirac matrices, Z is the nucleus charge, N is the number of electrons, $N-1$ is the number of core electrons, and ε_k, ψ_k are the energies and wave-functions of core orbitals. The orbitals of unpaired electron are generated by eqs (1, 2) in the potential of the frozen core.

The solution of eqs (1—3) gives the wave-functions of zero approximation. The exact Hamiltonian of an atom is of the form

$$H = \sum_{i=1}^{N} H_0(r_i) + U, \tag{4}$$

$$U = \sum_{i<j} \frac{e^2}{|\vec{r}_i - \vec{r}_j|} - \sum_{i=1}^{N} V^{(N-1)}(r_i). \tag{5}$$

The correlation corrections to the zero approximation we will find using the perturbation theory in residual Coulomb interaction U.

Parity non-conserving weak interaction of an electron with the nucleus looks as follows

$$h_{\mathrm{PNC}} = -\frac{G}{2\sqrt{2}}\,\varrho(r)\,Q_W\,\gamma_5, \tag{6}$$

where G is the Fermi constant, γ_5 is the Dirac matrix and $\varrho(r)$ is the nucleus density normalized by condition $\int \varrho(r)\,d^3r = 1$. We use the standard parametrization [17, 18]

$$\varrho(r) \infty \frac{1}{e^{(r-r_N)/D} + 1}. \tag{7}$$

Here $r_N = 1.105A^{1/3}$ fm, $D = 0.57$ fm, A is the mass number of the nucleus, Q_W is the weak charge of the nucleus. In the Weinberg—Salam theory $Q_W = -[N + Z(4\sin^2\theta_W - 1)]$, N is the number of neutrons, and Z is the number of protons. With the radiative correction taken into account [19]

$$Q_W = -[0.974\,N + Z(3.908\,\sin^2\theta_W(M_W) - 0.974)], \tag{8}$$

where $\sin^2\theta_W(M_W)$ is normalized at M_W.

The Hamiltonian h_{PNC} is one electron operator. Therefore it is convenient to insert it into Hartree—Fock equations. To do this one should replace in eqs (1—3) the orbital ψ with fixed parity to the orbital $\tilde{\psi}$ with mixed parity, and change the Hamiltonian

$$H_0(\psi) \to \tilde{H}_0(\tilde{\psi}) = H_0(\tilde{\psi}) + h_{\mathrm{PNC}}. \tag{9}$$

Let us expand the $\tilde{\psi}$ in the form $\tilde{\psi} = \psi + \delta\psi$, where $\delta\psi$ is the admixture with the same angular momentum but of the opposite parity. The equations for $\delta\psi$ can be derived from eqs (1—3)

$$(H_0 - \varepsilon_k)\, \delta\psi_k = -h_{\mathrm{PNC}}\, \psi_k - \delta V_{\mathrm{PNC}}\, \psi_k, \tag{10}$$

$$\delta V_{\mathrm{PNC}} = V^{(N-1)}(\tilde{\psi}_1, \ldots, \tilde{\psi}_{N-1}) - V^{(N-1)}(\psi_1, \ldots, \psi_{N-1}). \tag{11}$$

These equations should be solved in self-consistent way for the orbitals of the closed shells. The functions $\delta\psi$ for the states of unpaired electron are the solutions of eq. (10) at frozen core field $V^{(N-1)}$, δV_{PNC}.

Similar to the weak interaction the interaction with the external electromagnetic wave is the one electron operator. Therefore it is convenient to include it into Hartree—Fock equations as well. More precisely it is time-dependent Hartree—Fock (TDHF) method. Let us remind the reader the idea of this method. Let H_{int} be the Hamiltonian of the interaction of the electron with the external wave

$$H_{int} = \hat{f} e^{-i\omega t} + \hat{f}^{+} e^{i\omega t}. \tag{12}$$

Let also $\tilde{H}$ be the exact Hamiltonian of an atom

$$\tilde{H} = H + \sum_{i=1}^{N} h_{\mathrm{PNC}}(r_i). \tag{13}$$

Total TDHF function of an atom $\tilde{\Phi}$ is the Slatter determinant constructed from the single-particle orbitals $\tilde{\varphi}_{time\ dep}$

$$e^{-i\varepsilon_k t}\, \tilde{\varphi}_{td}^{(k)} = (\tilde{\psi}_k + \tilde{\chi}_k e^{-i\omega t} + \tilde{y}_k e^{i\omega t})\, e^{-i\varepsilon_k t}. \tag{14}$$

The tilde above the letters shows that the parity violation in corresponding functions is taken into account: $\tilde{\psi} = \psi + \delta\psi$, $\tilde{\chi} = \chi + \delta\chi$, $\tilde{y} = y + \delta y$. The equations for $\tilde{\chi}$, $\tilde{y}$ can be derived from condition

$$\delta \langle \tilde{\Phi} | \tilde{H} + H_{int} | \tilde{\Phi} \rangle = 0 \tag{15}$$

when varing in $\tilde{\chi}_k$, $\tilde{y}_k$

$$(\tilde{H}_0 - \varepsilon_k - \omega)\, \tilde{\chi}_k = -(\hat{f} + \delta\tilde{V}_f)\, \tilde{\psi}_k,$$
$$(\tilde{H}_0 - \varepsilon_k + \omega)\, \tilde{y}_k = -(\hat{f}^{+} + \delta\tilde{V}_f^{+})\, \tilde{\psi}_k, \tag{16}$$

$$\delta\tilde{V}_f\tilde{\Phi} = \sum_{k=1}^{N-1} \int d^3r_2\, \frac{e^2}{r_{12}} \{[\tilde{y}_k^{+}(2)\, \tilde{\psi}_k(2) + \tilde{\psi}_k^{+}(2)\, \tilde{\chi}_k(2)]\, \tilde{\Phi}(1) -$$

$$-[\tilde{y}_k^{+}(2)\, \tilde{\psi}_k(1) + \tilde{\psi}_k^{+}(2)\, \tilde{\chi}_k(1)]\, \tilde{\Phi}(2)\}. \tag{17}$$

Similar to the stationary case these equations should be solved self-consistently only for the $N-1$ core electrons ($V^{(N-1)}$-approximation). The wave-function of the external electron then can be found in frozen core field. This way allows us to find complete ortonormal set of the orbitals $\tilde{\psi}_{td}^{(k)}$ with quasienergy ε_k. At the first step we should solve the eqs (16, 17) for the quantities χ, y without parity violation. At the second step we can find the corrections $\delta\chi$, δy induced by the weak interaction. The equations for $\delta\chi$, δy can be obtained by decomposition of eqs (16, 17) in $\delta\chi$, δy, $\delta\varphi$, δH_0, δV_f. We do not present them here because the derivation is straightforward but the equations are combersome.

Let us consider transition of the external electron from the state $|\alpha\rangle$ to state $|\beta\rangle$ induced by the electric field (12). Expression for the transition amplitude $M_{\beta\alpha}$ can be obtained comparing the expression (14) with the formula of time-dependent perturbation theory

$$\tilde{\varphi}_{time\ dep}^{(\alpha)} = \tilde{\psi}_\alpha + \frac{M_{\beta\alpha}}{E_\alpha - E_\beta + \omega}\, \tilde{\psi}_\beta e^{-i\omega t}. \tag{18}$$

We keep here only the resonant term supposing that $\omega \sim E_\beta - E_\alpha$. From the other hand we see from eq. (16) that

$$\langle \tilde{\psi}_\beta | \tilde{\varphi}_{time\ dep}^{(\alpha)} \rangle \approx \langle \tilde{\psi}_\beta | \tilde{\chi}_\alpha e^{-u\omega t} \rangle = \frac{\langle \tilde{\psi}_\beta | \hat{f} + \delta\tilde{V}_f | \tilde{\psi}_\alpha \rangle}{E_\alpha - E_\beta + \omega}\, e^{-i\omega t}. \tag{19}$$

Comparing the equations (18) and (19) we get

$$M_{\beta\alpha} = \langle \tilde{\psi}_\beta | \hat{f} + \delta\tilde{V}_f | \tilde{\psi}_\alpha \rangle. \tag{20}$$

This formula corresponds to the well known random phase approximation (see e. g. Ref. [20]).

The Hamiltonian of interaction of an electron with electromagnetic wave depends on the gauge. Let φ be the scalar potential and $\vec{A}$ be the vector one. The conventional «length» form of interaction correspods to the choice

$$\varphi = -2\vec{E}_0 \vec{r} \cos(\omega t - \vec{k}\vec{r}).$$

$$\vec{A} = -\frac{2(\vec{E}_0 \vec{r})}{\omega} \cos(\omega t - \vec{k}\vec{r}), \qquad (21)$$

$$H_{int}^{(l)} = -e\,\vec{r}\,\vec{E}(t) = f_l e^{-i\omega t} + f_l^+ e^{i\omega t}, \qquad f_l = |e|\,\vec{r}\,\vec{E}_0.$$

Here $\vec{k}$ is a wave vector. We neglect the corrections $\sim kr$ and the magnetic interaction. Interaction in the «velocity» form corresponds to

$$\varphi = 0, \quad \vec{A} = -\frac{2\vec{E}_0}{\omega} \sin(\omega t - \vec{k}\vec{r}),$$

$$H_{int}^{(v)} = e\,\vec{\alpha}\,\vec{A}(t) = f_v e^{-u\omega t} + f_v^+ e^{i\omega t}, \qquad f_v = -i\,|e|\,\frac{\vec{\alpha}\,\vec{E}_0}{\omega}. \qquad (22)$$

The gauge (22) can be obtained from (21) by means of gauge transformation

$$\varphi \to \varphi + \frac{\partial}{\partial t}\theta, \quad \vec{A} \to \vec{A} - \vec{\nabla}\theta,$$

$$\theta = \frac{2\vec{E}_0 \vec{r}}{\omega} \sin(\omega t - \vec{k}\vec{r}). \qquad (23)$$

It is well known that in TDHF calculations amplitude (20) is gauge invariant [20, 21]. This statement is completely evident. The gauge transformation transfers single-determinant TDHF wave-function to the single-determinant one

$$\tilde{\Phi}' = \tilde{\Phi}\ \exp\left(i \sum_k \theta(r_k)\right). \qquad (24)$$

Therefore if $\tilde{\Phi}$ is the solution of eq. (15) in fixed gauge then $\tilde{\Phi}'$ is the solution of the same equation in other gauge. The relation between solutions of eqs (16) in different gauges can be found from formulae (23, 24). In linear approximation in external field

$$\tilde{\chi}_v = \tilde{\chi}_l - \frac{|e|\,\vec{r}\vec{E}_0}{\omega}\,\tilde{\Psi},$$

$$\tilde{y}_v = \tilde{y}_l + \frac{|e|\,\vec{r}\vec{E}_0}{\omega}\,\tilde{\Psi}. \qquad (25)$$

These relations has been obtained in Ref. [22]. Gauge invariance of the amplitude immediately follows from (25). According to (19) the amplitude equals to the residue of overlapping $\langle\tilde{\Psi}_\beta|\chi_\alpha\rangle$ at

Hartree—Fock frequency ($\omega = E_\beta - E_\alpha$), but gauge term in (25) has no pole.

Thus in TDHF the control of identity of the amplitudes in l- and v-form is only the control of numerical accuracy. The results of TDHF calculation of PNC E1-amplitudes for $6p_{1/2} \to 7p_{1/2}$ and $6p_{1/2} \to 6p_{3/2}$ transitions in ^{205}Tl and for $6s \to 7s$ transition in ^{133}Cs are presented in Table 1. The result for Cs coincides with our previous calculation [16] and with the result of Ref. [22]. (When comparing the amplitudes in v-form one should bear in mind that in Ref. [16] the calculation is carried out at experimental frequency but not at the Hartree—Fock one.) The TDHF-amplitude for $6p_{1/2} \to 7p_{1/2}$ transition practically precisely coincides with the results of Refs [11, 12].

3. CORRELATION CORRECTIONS TO PNC E1-AMPLITUDES

Let us consider the correlation corrections to TDHF amplitudes. There are the corrections of two types. Correction of the first type is given by the graphs with the photon radiation from the external line (Fig. 1). The dotted line corresponds to the photon, the circle—to the self-energy operator. The diagrams for the self-energy are presented at the Fig. 2. The wave line denotes the matrix element of the residual Coulomb interaction (5). The second type correlation correction corresponds to the photon radiation from the vertex (structural radiation). It is shown schematically at the Fig. 3. The photon attached to the circle means that it should be attached to the each internal electron line of the self-energy operator.

We begin from the consideration of the correction of the first type which is most important. The amplitude corresponding to the Fig. 1,*a* equals to

$$\sum_{\tilde{\gamma} \neq \tilde{\alpha}} \frac{\langle\tilde{\beta}|f + \delta\tilde{V}_f|\tilde{\gamma}\rangle\langle\tilde{\gamma}|\tilde{\Sigma}|\tilde{\alpha}\rangle}{\varepsilon_\alpha - \varepsilon_\gamma} = \sum_{\tilde{\gamma} \neq \tilde{\alpha}} \frac{\langle\tilde{\beta}|f + \delta\tilde{V}_f|\tilde{\gamma}\rangle\langle\tilde{\gamma}|\tilde{\Sigma}|\tilde{\alpha}\rangle}{\varepsilon_\beta - \varepsilon_\gamma - \omega} = \langle\tilde{y}_\beta|\tilde{\Sigma}|\tilde{\alpha}\rangle. \qquad (26)$$

We suppose that $\varepsilon_\beta > \varepsilon_\alpha$, ($\omega = \varepsilon_\beta - \varepsilon_\alpha$). Let us stress that the external field f we write down here with the accounting of shielding correction $\delta\tilde{V}_f$. The sum in formula (26) is carried out over the states $\tilde{\gamma} \neq \tilde{\alpha}$ and therefore the correction $\tilde{y}_\beta$ should satisfy the condition $\langle\tilde{y}_\beta|\tilde{\alpha}\rangle = 0$. With this additional condition the eqs (16) can be resolved despite of resonant frequency. At the same way one can write

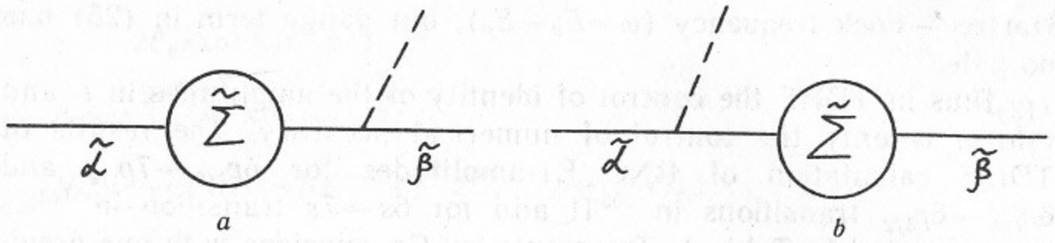

Fig. 1.

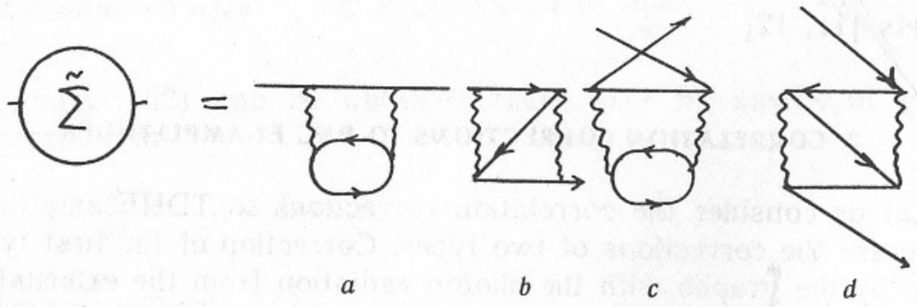

Fig. 2.

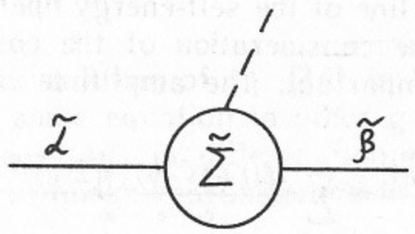

Fig. 3.

down the correction corresponding to the Fig. 1,*b*. Let us denote the total correction corresponding to Fig. 1 by $\Delta_{corr}(\tilde{\Sigma})$, then

$$\Delta_{corr}(\tilde{\Sigma}) = \langle \tilde{y}_\beta(\omega) | \tilde{\Sigma} | \tilde{\alpha} \rangle + \langle \tilde{\beta} | \tilde{\Sigma} | \tilde{\chi}_\alpha(\omega) \rangle, \tag{27}$$

where functions $\tilde{y}_\beta$ and $\tilde{\chi}_\alpha$ satisfy the conditions

$$\langle \tilde{y}_\beta(\omega) | \tilde{\alpha} \rangle = \langle \tilde{\beta} | \tilde{\chi}_\alpha(\omega) \rangle = 0, \qquad \omega = \varepsilon_\beta - \varepsilon_\alpha . \tag{28}$$

If we expand all the functions in linear approximation in weak interaction: $\tilde{y} = y + \delta y$, $\tilde{\Sigma} = \Sigma + \delta\Sigma$, ..., then we can represent the $\Delta_{corr}(\tilde{\Sigma})$ in the following form

$$\Delta_{corr}(\tilde{\Sigma}) = \Delta_{corr}(\Sigma) + \Delta_{corr}(\delta\Sigma), \tag{29}$$

where

$$\Delta_{corr}(\Sigma) = \langle \beta | \Sigma | \delta\chi_\alpha \rangle + \langle \delta\beta | \Sigma | \chi_\alpha \rangle + \langle \delta y_\beta | \Sigma | \alpha \rangle + \langle y_\beta | \Sigma | \delta\alpha \rangle, \tag{30}$$

$$\Delta_{corr}(\delta\Sigma) = \langle y_\beta | \delta\Sigma | \alpha \rangle + \langle \beta | \delta\Sigma | \chi_\alpha \rangle. \tag{31}$$

Let us repeat once more that the functions y_β, χ_α and PNC corrections δy_β, $\delta\chi_\alpha$ are the solutions of eqs (16) at additional conditions which follow from (28)

$$\langle \delta y_\beta | \alpha \rangle + \langle y_\beta | \delta\alpha \rangle = 0,$$

$$\langle \beta | \delta\chi_\alpha \rangle + \langle \delta\beta | \chi_\alpha \rangle = 0. \tag{32}$$

It is useful to estimate the relative contributions of different terms into the correlation correction before to present the results of numerical calculations. The quantity $\Delta_{corr}(\tilde{\Sigma})$ corresponds to the radiation from the external line. Therefore due to the small excitation energy of an external electron (the denominator in formula (26)) this contribution is enhanced by a factor $\Delta E_{int}/\Delta E_{ext}$ in comparison with the structural radiation (Fig. 3). Here ΔE_{int} is the typical excitation energy of the electrons from closed shells (denominator for the diagrams at the Fig. 3), ΔE_{ext} is the typical excitation energy of the external electron. Inside the correction $\Delta_{corr}(\tilde{\Sigma})$ the term $\Delta_{corr}(\Sigma)$ (formula(30)) is due to the weak interaction at the external line, the term $\Delta_{corr}(\delta\Sigma)$ (formula(31)) is due to the weak interaction inside the self-energy operator. Therefore here we have the same situation, the value of $\Delta_{corr}(\Sigma)$ is enhanced by a factor $\Delta E_{int}/\Delta E_{ext}$ in comparison with the $\Delta_{corr}(\delta\Sigma)$. We have used this hierarchy in Cs [16] where we have calculated only the dominating

contribution (30) using the method of Brueckner orbitals. For Tl the parameter $\Delta E_{int}/\Delta E_{ext}$ is not so large as for Cs and therefore here we take into account all the second order correlation corrections.

The results of the l-form calculation of the contributions into the correlation correction $\Delta_{corr}(\tilde{\Sigma})$ (see eqs (30, 31)) are presented in the Table 2. The results of the calculation in the v-form are presented in the Table 3. We want to remind the reader that we suppose that transition goes from state $|\alpha\rangle$ to state $|\beta\rangle$, and $\varepsilon_\beta > \varepsilon_\alpha$. The self-energy operator we calculate by direct summation over the intermediate states (see Ref. [16]).

In v-form the operator f depends on frequency (see eq. (22)), and the correlations shift the frequency

$$\omega = \varepsilon_\beta - \varepsilon_\alpha \to \omega + \Delta\omega,$$

$$\Delta\omega = \langle\beta|\,\Sigma|\,\beta\rangle - \langle\alpha|\,\Sigma|\,\alpha\rangle. \qquad (33)$$

So the transition amplitude

$$\langle\tilde{\beta}| - \frac{i\vec{\alpha}}{\omega}\,|\tilde{\alpha}\rangle \to \langle\tilde{\beta}| - \frac{i\vec{\alpha}}{\omega+\Delta\omega}\,|\tilde{\alpha}\rangle \approx$$

$$\approx \langle\tilde{\beta}| - \frac{i\vec{\alpha}}{\omega}\,|\tilde{\alpha}\rangle - \frac{\Delta\omega}{\omega}\langle\tilde{\beta}| - \frac{i\vec{\alpha}}{\omega}\,|\tilde{\alpha}\rangle \approx$$

$$\approx \langle\tilde{\beta}| - \frac{i\vec{\alpha}}{\omega}\,|\tilde{\alpha}\rangle - \frac{\Delta\omega}{\omega}\langle\tilde{\beta}|\vec{r}|\tilde{\alpha}\rangle. \qquad (34)$$

Thus due to dependence of f_v on frequency in v-form the additional correlation correction arises

$$-\frac{\Delta\omega}{\omega}\langle\tilde{\beta}|\vec{r}|\,\alpha\rangle = -\frac{\Delta\omega}{\omega}[\langle\delta\beta|\vec{r}|\,\alpha\rangle + \langle\beta|\vec{r}|\,\delta\alpha\rangle]. \qquad (35)$$

The numerical values of this correction are given in the last line of the Table 3. In the formula (35) we write down the matrix element $\langle\tilde{\beta}|\vec{r}|\,\tilde{\alpha}\rangle$ without the polarization correction. The polarization effect in correlation correction is at least the third order term of perturbation theory. Therefore strictly speaking from derivation (34) we can not conclude should we take into account the polarization in formula (35) or we should not. However we want to claim the gauge invariance of correlation correction. Then transforming the formulae (27) and (28) from l- to v-form by means of (25) it is easy to ve-

rify that precisely the formula (35) without accounting of polarization is correct.

In accordance with our preliminary estimation in l-form the value of $\Delta_{corr}(\delta\Sigma)$ is suppressed in comparison with the $\Delta_{corr}(\Sigma)$ but for the transition $6p_{1/2} \to 6p_{3/2}$ in the Tl this is suppression only by a factor 2 (see Table 2). As far as the correlations in the v-form are concerned the real suppression takes place only for the transition $6s \to 7s$ in Cs (see Table 3). One can easily understand this fact. The matrix elements of electromagnetic interaction in v-form are large for high frequency virtual transitions and therefore the parameter $\Delta E_{int}/\Delta E_{ext}$ for Tl is of the order of unit. From the Table 3 we see that for the transition $6p_{1/2} \to 6p_{3/2}$ in Tl the contributions to the correlation correction in the v-form are of the same order as the main Hartree—Fock amplitude (Table 1). Due to this fact the calculation in v-form is rather unstable. We will return to this question after the discussion of the structural radiation.

Structural radiation has two parts. In the first one the weak interaction is included into the external lines and the electromagnetic

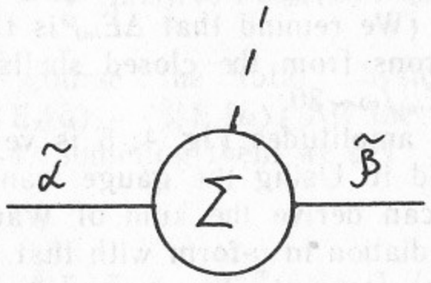

Fig. 4.

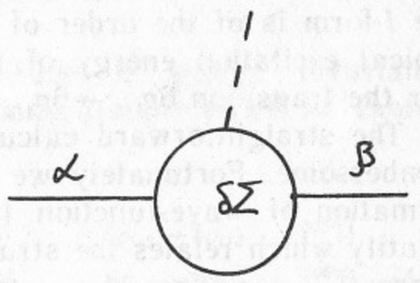

Fig. 5.

vertex itself is parity conserving (Fig. 4). In the second part the weak interaction is included into the electromagnetic vertex (Fig. 5). Let us discuss at first the electromagnetic interaction in the l-form. In this case the structural radiation is small and we can limit ourselves by very rough numerical calculation. The results of this calculation for diagram Fig. 4 are presented in first column of Table 4.

The amplitude corresponding to Fig. 5 for $6p_{1/2} \to 7p_{1/2}$ transition in Tl and for $6s \to 7s$ transition in Cs is very small and therefore we estimate it only for the transition $6p_{1/2} \to 6p_{3/2}$ in Tl. The estimation is very simple. Here the diagrams with the intermediate $6s$- and $6p$-states dominates. An example of such diagram is presented at Fig. 6. The cross denotes the PNC weak interaction. The number of

these graphs is rather large, but it is not so complicated to calculate all of them. The result is given in the second column of Table 4.

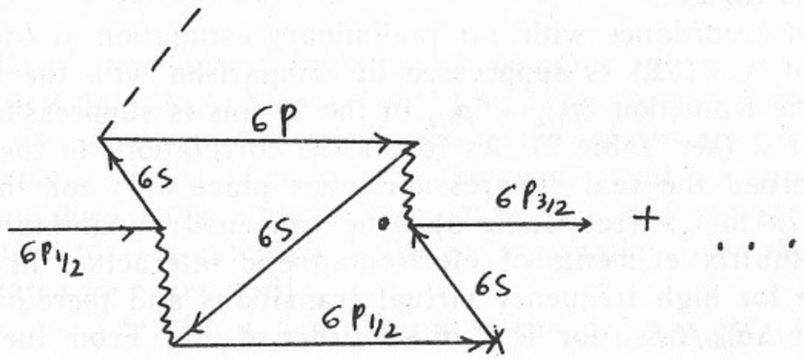

Fig. 6.

In the v-form the structural radiation is not small and therefore we need more accurate calculation than for the l-form. The reason of enhancement of structural radiation in v-form is evident. The operator of electromagnetic interaction in v-form is proportional to $1/\omega$ (eq. (22)). Therefore the factor of enhancement with respect to the l-form is of the order of $\Delta E_{int}/\omega$. (We remind that ΔE_{int} is the typical excitation energy of the electrons from the closed shells.) For the transition $6p_{1/2} \to 6p_{3/2}$ in Tl $\Delta E_{int}/\omega \sim 30$.

The straightforward calculation of amplitudes Fig. 4, 5 is very cumbersome. Fortunately we can avoid it. Using the gauge transformation of wave-function (25) we can derive the kind of Ward identity which relates the structural radiation in v-form with that in l-form

$$\langle\tilde{\beta}|\,\vec{\Gamma}^{(v)}_{str}\,|\tilde{\alpha}\rangle = \langle\tilde{\beta}|\,\vec{\Gamma}^{(l)}_{str}\,|\tilde{\alpha}\rangle + \frac{e}{\omega}\langle\tilde{\beta}|\,\tilde{\Sigma}(\varepsilon_\beta)\,\vec{r} - \vec{r}\,\tilde{\Sigma}(\varepsilon_\alpha)\,|\tilde{\alpha}\rangle. \tag{36}$$

Here $\Gamma_{str} = \Gamma(\text{Fig. 4}) + \Gamma(\text{Fig. 5})$. One can separates the eq. (36) into the parts corresponding Fig. 4 and Fig. 5.

$$\langle\tilde{\beta}|\,\vec{\Gamma}^{(v)}_{\text{Fig.4}}\,|\tilde{\alpha}\rangle = \langle\tilde{\beta}|\,\vec{\Gamma}^{(l)}_{\text{Fig.4}}\,|\tilde{\alpha}\rangle + \frac{e}{\omega}\langle\tilde{\beta}|\,\Sigma(\varepsilon_\beta)\,\vec{r} - \vec{r}\,\Sigma(\varepsilon_\alpha)\,|\tilde{\alpha}\rangle,$$

$$\langle\beta|\,\vec{\Gamma}^{(v)}_{\text{Fig.5}}\,|\alpha\rangle = \langle\beta|\,\vec{\Gamma}^{(l)}_{\text{Fig.5}}\,|\alpha\rangle + \frac{e}{\omega}\langle\beta|\,\delta\Sigma(\varepsilon_\beta)\,\vec{r} - \vec{r}\,\delta\Sigma(\varepsilon_\alpha)\,|\alpha\rangle. \tag{37}$$

Using this formulae we can easy calculate the structural radiation in v-form. The results of calculation are presented in Table 4.

In the second order of perturbation theory in residual interaction

there is **one** else term which is due to the decrease of normalization of the single particle orbitals $|\alpha\rangle$, $|\beta\rangle$ because of the admixture of the many-particle states. It has the form

$$\Delta_{norm} = \frac{1}{2}\langle\tilde{\beta}|\,f + \delta\tilde{V}_f\,|\tilde{\alpha}\rangle\left[\langle\beta|\,\frac{\partial\Sigma}{\partial E}\,|\beta\rangle + \langle\alpha|\,\frac{\partial\Sigma}{\partial E}\,|\alpha\rangle\right]. \tag{38}$$

This correction is proportional to the TDHF amplitude $\langle\tilde{\beta}|\,f + \delta\tilde{V}_f\,|\tilde{\alpha}\rangle$ and therefore it is gauge invariant by itself. The values of this correction are given in the last column of Table 4.

Thus we have calculated all the corrections of the second order in residual Coulomb interaction. In the l-form the final result for the PNC E1-amplitude looks as follows

$$\langle\tilde{\beta}|\,E_z\,|\tilde{\alpha}\rangle = \langle\tilde{\beta}|\,E_z\,|\tilde{\alpha}\rangle_{\text{TDHF}} + \Delta^{(l)}_{corr}(\Sigma) + \Delta^{(l)}_{corr}(\delta\Sigma) + \Gamma^{(l)}_{str} + \Delta_{norm}. \tag{39}$$

The same amplitude in the v-form

$$\langle\tilde{\beta}|\,E_z\,|\tilde{\alpha}\rangle = \langle\tilde{\beta}|\,E_z\,|\tilde{\alpha}\rangle_{\text{TDHF}} -$$

$$-\frac{\Delta\omega}{\omega}\langle\tilde{\beta}|\,er_z\,|\tilde{\alpha}\rangle + \Delta^{(v)}_{corr}(\Sigma) + \Delta^{(v)}_{corr}(\delta\Sigma) + \Gamma^{(v)}_{str} + \Delta_{norm}. \tag{40}$$

Of course the total physical amplitude is gauge invariant: $\langle\tilde{\beta}|\,E_z\,|\tilde{\alpha}\rangle_l = \langle\tilde{\beta}|\,E_z\,|\tilde{\alpha}\rangle_v$ All the contributions are presented in Tables 1—4. Summing them we get

$$\langle\tilde{\beta}|\,E_z\,|\tilde{\alpha}\rangle = -10^{-10}(i|e|\,a_B\,Q_W/-N)\left\{\begin{array}{lll} 0.81 & 6p_{1/2}\to 7p_{1/2} & {}^{205}\text{Tl} \\ 2.65 & 6p_{1/2}\to 6p_{3/2} & {}^{205}\text{Tl} \\ -0.0897 & 6s\to 7s & {}^{133}\text{Cs} \end{array}\right. \tag{41}$$

This is the result of consistent calculation. In the next section we will slightly correct these values basing on experimental data on energy levels, oscillator strengths, hyperfine structure, and on the theoretical estimation of the third order correlation correction.

In conclusion of this section we would like to comment on the v-form calculation. As stated above for $6p_{1/2} \to 6p_{3/2}$ transition in Tl in the v-form the contributions to the correlation correction (Tables 3, 4) were of the same order of magnitude as the main TDHF-amplitude (Table 1). Only due to strong numerical compensations the total correlation correction is small. The reason of compensation was pointed above. In v-form each contribution into correlation cor-

rection is proportional to $1/\omega$ and therefore it is large at small ω. At the same time the total correction is independent of ω. Thus the calculation in the v-form is unstable.

4. THE FINAL RESULTS. ANALYSIS OF ACCURACY OF CALCULATION

The most reliable way to estimate the theoretical accuracy is calculation of the quantities which are known from experiment. In Ref. [23] we have calculated the energy levels, hyperfine constants and the amplitudes of the allowed E1-transitions. All these values are very important for the semiempirical calculation of the PNC E1-amplitudes. The results obtained in Ref. [23] are as follow. The energy level (ionization potential) is reproduced with accuracy:

$6p$ — 2.9%, $7p$ — 0.7%, $7s$ — 0.7%.

The accuracy of fine structure calculation is:

$6p$ — 2%, $7p$ — —2.5%.

The calculation of hyperfine structure of s and $p_{1/2}$ states controls the electron wave-functions at the nucleus. Here we have the accuracy:

$6p_{1/2}$ — —2.4%, $7p_{1/2}$ — (-4 ± 3)%, $7s$ — —0,5%.

The uncertainty in the value for $7p$-level is due to the experimental error.

In the Ref. [23] for allowed E1-transitions we estimated only the dominating diagram in the structural radiation. In the present work we recalculate this radiation taking into account all the second order diagrams. The new results for E1-amplitudes are presented in the Table 5. Of course they almost precisely coincide with the old results of Ref. [23]. In the semiempirical calculation of the PNC in $6p_{1/2}\to6p_{3/2}$ transition [13] the $6p_{3/2}\to7s$ E1-amplitude is the most important. For it we have the accuracy:

E1$(6p_{3/2}\to7s)$ — (-2.8 ± 2.1)%.

In semiempirical calculation of PNC E1-amplitude for $6p_{1/2}\to6p_{3/2}$ transition all the intermediate states contribute with the

positive sign [13]. Since there is no compensation we think that the relative accuracy of our calculation of this amplitude is the same as the accuracy of calculation of the quantities considered above, namely 3%. The position of the central point in the final result can be refined basing on following observation. Usually we slightly overestimate the magnitude of the correlation correction. It means that the diagrams of higher orders compensate the part of the second order contribution. To take this fact into account we can multiply the second order correlation correction by a factor which is less than unit. If we introduce the common factor 0.85 then we have the accuracy:

energy of $6p_{1/2}$ — 0.3%,

energy of $7s$ — —0.5%,

fine structure of $6p$ — 0.6%,

hyperfine structure of $6p_{1/2}$ — —2.1%,

hyperfine structure of $7s$ — —2.9%,

E1$(6p_{1/2}\to7s)$ amplitude — (-0.8 ± 2.1)%.

We can go further, namely we can introduce the different factors for different waves (0.95 for Σ_s and 0.8 for $\Sigma_{p_{1/2}}$). In this case the agreement with experiment for every quantity is better than 1%.

For the PNC E1$(6p_{1/2}\to6p_{3/2})$ amplitude the both ways (common factor and different factors) give the result

$$\langle 6p_{3/2}\tfrac{1}{2}\,|\,E_z\,|\,6p_{1/2}\tfrac{1}{2}\rangle = -2.70\cdot10^{-10}\cdot(i\,|\,e|\,a_B Q_W/-N).$$

Besides the comparison with experiment we have pure theoretical estimation of accuracy from the analysis of higher orders contribution. First of all we have taken into account the iterations of self energy correction by the method of Brueckner orbitals, i. e. we have introduced the Σ into the equations for orbitals (see e. g. Ref. [16]). In linear approximation in Σ the «Brueckner» calculation is identical to the «straightforward» one. Due to the nonlinear in Σ terms the «Brueckner» result is slightly larger then the «straigtforward» one (2.74 instead of 2.65). However the «Brueckner» value exceeds the refine one (2.70) only by a 1.5%. From our view important diagrams of the third order in residual interaction correspond to accounting of self energy corrections in the core polarization effect (including the shift of the transition frequency). Our esti-

mation for the contribution of these diagrams lies in the range 0.4—2.0%.

The error which is due to inaccuracy of neutron and proton distributions in the nucleus is less then 1%.

Thus we have two ways for estimation of accuracy of our calculation. The first one is the comparison with experiment of the quantities which are known from experiment. The second way is the estimation of higher order corrections of perturbation theory. The both these ways give the value of the theoretical error $\lesssim 3\%$. The final result for the PNC E1-amplitude of transition $6p_{1/2} \to 6p_{3/2}$ in ^{205}Tl is

$$\langle 6p_{3/2}\frac{1}{2}|E_z|6p_{1/2}\frac{1}{2}\rangle = -2.70\cdot(1\pm 0.03)\cdot 10^{-10}\cdot(i|e|a_B Q_W/-N). \quad (42)$$

The calculation of PNC for the transition $6p_{1/2} \to 7p_{1/2}$ is less stable. In semiempirical calculation there are strong compensations between contributions of different intermediate states. Therefore the inaccuracy can be larger then in the calculation of the energy levels, hyperfine structure and the allowed E1-amplitudes. For the $6p_{1/2} \to 7p_{1/2}$ the result is more sensitive to the high order perturbation theory contributions. Accounting of nonlinear in Σ corrections by the Brueckner orbitals method increases the PNC-amplitude by 5%. From the other hand accounting of self-energy corrections in the core polarization effect decreases the result by 4%. In this situation we think that the value 6% is the reasonable estimation of accuracy. The final result looks as follows:

$$\langle 7p_{1/2}\frac{1}{2}|E_z|6p_{1/2}\frac{1}{2}\rangle = -0.81\cdot(1\pm 0.06)\cdot 10^{-10}\cdot(i|e|a_B Q_W/-N). \quad (43)$$

Let us note that the sensitivity of this amplitude to the self-energy correction in the core polarization is due to the closeness of $6p_{1/2} \to 7p_{1/2}$ transition frequency to the frequency of $6s \to 6p$ core excitation. Configurations corresponding to these excitations are $6s^2 7p_{1/2}$ and $6s6p^2$. The mixing of these configurations by the PNC weak interaction together with the residual Coulomb interaction gives contribution to the PNC E1-amplitude. The experimental interval between the level $6s^2 7p_{1/2}$ and the lowest level of configuration $6s6p^2$ is only 11060 cm^{-1}. It is very complicated to reproduce in calculation so small splitting. This splitting and hence the corresponding contribution to the PNC E1-amplitude are sensitive to any correction.

In conclusion we consider the $6s \to 7s$ transition in Caesium. Here the accuracy of calculation of energy levels [24], fine and hyperfine intervals [25], and allowed E1-amplitudes [16] is rather high. Therefore the introduction of correction factor for correlations does not influence practically on the PNC E1-amplitude. Accounting of nonlinear in Σ corrections by the Brueckner orbitals method decreases the result by 2%. Accounting of self-energy corrections in the core polarization effect changes the result less then by 1%. The final value of PNC E1-amplitude for ^{133}Cs

$$\langle 7s\frac{1}{2}|E_z|6s\frac{1}{2}\rangle = 0.89\cdot(1\pm 0.02)\cdot 10^{-11}\cdot(i|e|a_B Q_W/-N). \quad (44)$$

In analogy with Cs we have calculated the PNC E1-amplitude for $5s \to 6s$ transition in ^{85}Rb

$$\langle 6s\frac{1}{2}|E_z|5s\frac{1}{2}\rangle = 0.138\cdot 10^{-11}\cdot(i|e|a_B Q_W/-N). \quad (45)$$

For the Rubidium we have not calculated the hyperfine structure and the allowed E1-amplitudes. Therefore we have no the detailed analysis of accuracy. However we think that the accuracy for Rb is the same as for Cs, i. e. 2%.

Now we can compare our results for PNC E1-amplitudes with the experimental data and with the results of other calculations. In the Tables 6—8 these amplitudes are presented. Experimental data are recalculated for $\sin^2\theta_W = 0.22$. Let us note that the results of present work are very close to that obtained by the semiempirical method [7, 13, 26]. The reason of this fact is evident. In the effective potential method one fits the self-energy correlation correction by fitting the energy levels. We have seen in present work that precisely the self-energy correction is most important for PNC E1-amplitude. As far as the electromagnetic core polarization is concerned it is taken into account in semiempirical method by using the experimental E1-amplitudes.

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Table 1

The PNC E1-Amplitudes in l and v-forms
(Units $10^{-10}(i\|e\|a_B Q_W/-N)$)

Transition		HF	HF + Weak Polarization	HF + Weak Polarization + + Electromagnetic Polarization
^{205}Tl	l	−0.663	−0.991	−0.965
$\langle 7p_{1/2}\frac{1}{2}\|E_z\|6p_{1/2}\frac{1}{2}\rangle$	v	−0.937	−0.754	−0.965
^{205}Tl	l	−2.472	−3.978	−3.034
$\langle 6p_{3/2}\frac{1}{2}\|E_z\|6p_{1/2}\frac{1}{2}\rangle$	v	1.237	−3.873	−3.034
^{133}Cs	l	0.0742	0.0922	0.0887
$\langle 7s\frac{1}{2}\|E_z\|6s\frac{1}{2}\rangle$	v	0.0596	0.0826	0.0887

The first column corresponds to simple Hartree-Fock calculation. The second one is the calculation with the weak interaction included into the HF equations. The third column corresponds to the calculation with the both weak and electromagnetic interaction included into the HF equations.

Table 2

The Contributions in the l-form into the Correlation Correction $\Delta_{corr}(\tilde{\Sigma}) = \Delta_{corr}(\Sigma) + \Delta_{corr}(\delta\Sigma)$ Corresponding to Fig. 2.
(Units $10^{-10}(i\|e\|a_B Q_W/-N)$)

	^{205}Tl $\langle 7p_{1/2}\frac{1}{2}\|E_z\|6p_{1/2}\frac{1}{2}\rangle$	^{205}Tl $\langle 6p_{3/2}\frac{1}{2}\|E_z\|6p_{1/2}\frac{1}{2}\rangle$	^{133}Cs $\langle 7s\frac{1}{2}\|E_z\|6s\frac{1}{2}\rangle$
$\langle\beta\|\Sigma\|\delta\chi_\alpha\rangle$	−0.029	0.144	0.0064
$\langle\delta\beta\|\Sigma\|\chi_\alpha\rangle$	−0.037	0.022	0.0182
$\langle\delta y_\beta\|\Sigma\|\alpha\rangle$	0.364	0.239	−0.0090
$\langle y_\beta\|\Sigma\|\delta\alpha\rangle$	−0.131	0.337	−0.0149
$\Delta_{corr}(\Sigma)$	0.167	0.742	0.0007
$\langle y_\beta\|\delta\Sigma\|\alpha\rangle$	−0.012	−0.231	−0.0006
$\langle\beta\|\delta\Sigma\|\chi_\alpha\rangle$	0.010	−0.100	0.0012
$\Delta_{corr}(\delta\Sigma)$	−0.002	−0.331	0.0006

Table 3

The Contributions in the v-form into the Correlation Correction $\Delta_{corr}(\tilde{\Sigma}) = \Delta_{corr}(\Sigma) + \Delta_{corr}(\delta\Sigma)$ Corresponding to Fig. 2 (Units $10^{-10}(i|e|a_B Q_W/-N)$)

	^{205}Tl $\langle 7p_{1/2}\frac{1}{2}\|E_z\|6p_{1/2}\frac{1}{2}\rangle$	^{205}Tl $\langle 6p_{3/2}\frac{1}{2}\|E_z\|6p_{1/2}\frac{1}{2}\rangle$	^{133}Cs $\langle 7s\frac{1}{2}\|E_z\|6s\frac{1}{2}\rangle$
$\langle\beta\|\Sigma\|\delta\chi_\alpha\rangle$	0.034	0.759	0.0173
$\langle\delta\beta\|\Sigma\|\chi_\alpha\rangle$	−0.255	0.201	0.0081
$\langle\delta y_\beta\|\Sigma\|\alpha\rangle$	0.238	−3.547	0.0209
$\langle y_\beta\|\Sigma\|\delta\alpha\rangle$	−0.253	4.618	−0.0286
$\Delta_{corr}(\Sigma)$	−0.236	2.031	0.0177
$\langle y_\beta\|\delta\Sigma\|\alpha\rangle$	0.043	−1.335	−0.0009
$\langle\beta\|\delta\Sigma\|\chi_\alpha\rangle$	0.078	−1.753	0.0008
$\Delta_{corr}(\delta\Sigma)$	0.121	−3.088	−0.0001
$-\frac{\Delta\omega}{\omega}\langle\tilde{\beta}\|r_z\|\tilde{\alpha}\rangle$	0.198	0.363	−0.0157

In the last line the values of correction induced by the frequency shift are presented.

Table 4

The Structural Radiation in l and v-forms: $\Gamma_4 = \Gamma$(Fig. 4), $\Gamma_5 = \Gamma$(Fig. 5)

	l-form			v-form			Δ_{norm}
	Γ_4	Γ_5	$\Gamma_{str} = \Gamma_4 + \Gamma_5$	Γ_4	Γ_5	$\Gamma_{str} = \Gamma_4 + \Gamma_5$	
^{205}Tl $\langle 7p_{1/2}\frac{1}{2}\|E_z\|6p_{1/2}\frac{1}{2}\rangle$	−0.02	0.	−0.02	0.185	−0.123	0.062	0.013
^{205}Tl $\langle 6p_{3/2}\frac{1}{2}\|E_z\|6p_{1/2}\frac{1}{2}\rangle$	−0.06	−0.03	−0.09	−1.712	2.727	1.015	0.06
^{133}Cs $\langle 7s\frac{1}{2}\|E_z\|6s\frac{1}{2}\rangle$	0.0003	0.	0.0003	−0.001	0.0007	−0.0003	−0.0006

In the last column the normalization contribution is presented. Units $10^{-10}(i|e|a_B Q_W/-N)$.

Table 5

The Radial Integrals $R_{\beta\alpha}$ for Allowed E1-Amplitudes in Tl

$\langle\beta|E_z|\alpha\rangle = ea_B R_{\beta\alpha}\langle\beta|r_z/r|\alpha\rangle$

Transition	TDHF	TDHF + + correlations	Experiment Ref. [27]
$6p_{1/2}\to 7s$	2.32	2.11	2.23 (6)
$6p_{3/2}\to 7s$	3.13	2.75	2.83 (6)
$6p_{1/2}\to 8s$	0.75	0.67	0.67 (3)
$6p_{3/2}\to 8s$	0.70	0.67	—
$6p_{1/2}\to 6d_{3/2}$	−2.15	−2.07	−1.99 (8)
$6p_{3/2}\to 6d_{3/2}$	−2.94	−2.70	−2.64 (13)

Table 6

The Amplitude $\langle 7p_{1/2}\,\frac{1}{2}|E_z|6p_{1/2}\,\frac{1}{2}\rangle$ for ^{205}Tl in Units $10^{-10}(i|e|a_B Q_W/-N)$

Experiment	Drell and Commins [4] Tanner and Commins [5]	$-0.75(1\pm0.19)$
Calculation	Bouchiat and Bouchiat [6]	-0.66
	Sushkov et al. [7]	-0.80
	Neuffer and Commins [8]	$-0.83(1\pm0.22)$
	Das et al. [9]	$-0.683(1\pm0.05)$
	Plummer and Grant [10]	-0.998
	Johnson et al. [11]	-0.971
	Martensson-Pendrill [12]	-0.956
	Present work	$-0.81(1\pm0.06)$

Table 7

The Amplitude $\langle 6p_{3/2}\,\frac{1}{2}|E_z|6p_{1/2}\,\frac{1}{2}\rangle$ for ^{205}Tl in Units $10^{-10}(i|e|a_B Q_W/-N)$

Calculation	Novikov et al. [13]	$-2.87(1\pm0.2)$
	Henley and Wilets [14]	-3.87
	Henley et al. [15]	-2.55
	Neuffer and Commins [8]	-3.8
	Present work	$-2.70(1\pm0.03)$

Table 8

The Amplitude $\langle 7s\,\frac{1}{2}|E_z|6s\,\frac{1}{2}\rangle$ for ^{133}Cs in Units $10^{-10}(i|e|a_B Q_W/-N)$

Experiment	Bouchiat et al. [33]	$0.88(1\pm0.11)$
	Gilbert et al. [34]	$0.96(1\pm0.08)$
Calculation	Bouchiat and Bouchiat [6]	1.33
	Loving and Sandars [28]	1.15
	Neuffer and Commins [29]	1.00
	Bouchiat et al. [30]	0.97 ± 0.1
	Kuchiev et al. [31]	0.75
	Das et al. [32]	1.06
	Dzuba et al. [16]	0.88 ± 0.03
	Martensson-Pendrill [22]	0.886
	Johnson et al. [11]	0.890
	Schafer et al. [35]	0.92
	Bouchiat and Piketty [26]	$0.935\pm0.02\pm0.03$
	Present work	$0.89(1\pm0.02)$

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Relativistic Many-Body Calculation of Parity Non-Conservation in Thallium

В.А. Дзюба, В.В. Фламбаум, П.Г. Сильвестров, О.П. Сушков

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