Effects of nonconservation of parity in two-electron atoms and ions

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Effects of nonconservation of parity in two-electron atoms and ions arising in the presence of neutral weak currents are considered. It is shown that for an ion of the carbon isotope $C^{13}V$ the degree of asymmetry in the emission of photons with respect to the direction of the magnetic field H attains the value of unity for $H \sim 10^6$ G. In weak magnetic fields the degree of asymmetry of photon emission and the degree of circular polarization is of the order of 10^{-4} for He atoms and $C^{13}V$ and Cu XXVIII ions.

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The existence of a weak interaction between the electrons and the nucleus in an atom can become evident as a result of different phenomena of nonconservation of parity in atomic transitions involving radiation[1,2]. These phenomena include the appearance of circular polarization of the emitted photons and the asymmetry of emission of photons with respect to a magnetic field or to the initial spin of the atom. In the presence of weak neutral currents violation of parity in atoms occurs in terms of the order of $G\eta^2 \sim 10^{-16}$, where G is the Fermi weak interaction constant, $\eta = m\alpha Z$ is the characteristic momentum of an electron in an atom, m is the electron mass, α is the fine structure constant, Z is the nuclear charge. However, there exist different mechanisms which amplify the effects of nonconservation of parity[3-7]. In order for it to be possible to observe these effects it is necessary that the line being studied in a spectrum should be sufficiently intense, the degree of polarization or of asymmetry P should be not too small and the accidentally present electric or magnetic fields should not be able to simulate the effects of nonconservation of parity.

For one-electron or single-valent atoms or ions in all cases that have been discussed $P \leq 10^{-4}$, the lines being investigated are very weak, and the limitations imposed on the external accidental fields are very rigid[5]. We consider two-electron or doubly-valent atoms and ions and shall study effects of nonconservation of parity involving the metastable level 2^1S_0[7]. The one-quantum decay mode of the level 2^1S_0 into the ground state with the emission of M1 photons is determined by the hyperfine interaction V_{HF}, as a result of which the 2^3S_1 level is admixed to the 2^1S_0 level (the remaining levels of the same parity lie far away and one can neglect their admixture). Thus, this mode is absent for a spinless nucleus.

The expression for V_{HF} useful for the calculation of matrix elements with S-states has the form

$$V_{HF} = \frac{8\pi}{3} g\mu_0^2 \frac{m}{m_p} \sum_{j=1,2} \sigma_j \mathbf{T} \delta^{(3)}(\mathbf{r}_j), \tag{1}$$

where σ_j are the Pauli matrices, T is the nuclear spin, r_j are the electron coordinates, μ_0 is the Bohr magneton, m is the electron mass, m_p is the proton mass, g is the gyromagnetic ratio for the nucleus.

When parity is not conserved also another single-quantum mode appears for the same decay involving the emission of E1-photons. Violation of parity is due to the weak interaction V_W which mixes S- and P- levels of the atom and which is proportional to the product of the electron and the nuclear neutral currents. The expression for V_W can be obtained from the amplitude of the weak interaction process, which in the momentum representation has the form

$$2^{-1/2} G \bar{\psi}_e \gamma_\mu^e (g_e + h_e \gamma_5^e) \psi_e \bar{\psi}_{p,n} \gamma_\mu^{(p,n)} (g_{p,n} + h_{p,n} \gamma_5^{(p,n)}) \psi_{p,n}, \tag{2}$$

where ψ_e and $\psi_{p,n}$ are the electron and the nucleon (proton or neutron) wave functions, G is the Fermi constant, g and h are coefficients depending on the model being used, γ_μ are the Dirac matrices. We have omitted in the nuclear current in (2) terms containing q_μ and $q_\nu \delta_{\mu\nu}$, where q is the transferred momentum, $\delta_{\mu\nu} = -i(\gamma_\nu\gamma_\mu - \gamma_\mu\gamma_\nu)$. The smallness of these terms is related to the smallness of the characteristic atomic momenta compared to the characteristic nuclear momenta.

Going over to the nonrelativistic approximation for the nucleus and taking into account only that part of (2) which violates parity we obtain

$$2^{-1/2} G \{ h_e g_{p,n} (\psi_e^+ \gamma_5^e \psi_e)(\psi_{p,n}^+ \psi_{p,n}) + g_e h_{p,n} (\psi_e^+ \boldsymbol{\alpha}^e \psi_e)(\psi_{p,n}^+ \boldsymbol{\sigma}^{(p,n)} \psi_{p,n}) \},$$

where $\sigma^{(p,n)}$ are the Pauli matrices for the proton (neutron). Averaging over the motion of the nucleons in the nucleus we obtain an expression for V_W in the momentum representation:

$$V_W = 2^{-1/2} G (a\gamma_5 + b\boldsymbol{\alpha}\mathbf{T}),$$

where T is the nuclear spin, a and b are constants characterizing the magnitude of the neutral currents. In our case the term with γ_5 is not operative, since it does not mix states with different total electron angular momenta, such as 2^1S_0 and 2^3P_1. The constant b is determined by the relation

$$b\mathbf{T} = \left\langle \psi_T^n \left| \varkappa_p \sum_i \sigma_i^{(p)} + \varkappa_n \sum_i \sigma_i^{(n)} \right| \psi_T^n \right\rangle,$$

where ψ_T^n is the wave function describing the state of the nucleus with spin T, while the summations are extended over all the protons (neutrons) of the nucleus. In the Weinberg model[8] $\kappa_p = -\kappa_n = 0.24$.

In the nonrelativistic approximation in coordinate space the operator V_W for two electrons has the form

$$V_W = \frac{bG}{2^{3/2} m} \sum_{j=1,2} \{ \mathbf{T}\nabla_j + i[\mathbf{T}\sigma_j]\nabla_j \} \delta^{(3)}(\mathbf{r}_j), \tag{3}$$

where σ_j are the Pauli matrices for the electron.

The probability of a single quantum transition taking into account both decay modes in the presence of an external magnetic field H and also of a random electric field E is equal to

$$dw^{1\gamma} = w_{S_1}^{1\gamma} \{ 1 + \mathbf{nh}P + (\boldsymbol{\varepsilon}\mathbf{h})(\mathbf{nh})Q \} \frac{d\Omega}{4\pi}, \tag{4}$$

where $\mathbf{h} = \mathbf{H}/H$, $\epsilon = \mathbf{E}/E$, $\mathbf{n} = \mathbf{k}/k$, $\mathbf{k}$ is the photon momentum, $w_{S1}^{1\gamma}$ is the total probability of a single quantum transition in the absence of a weak interaction. In the above we have $w_{S1}^{1\gamma} = \mathscr{K}^2 w_{S3}^{1\gamma}$, where $w_{S3}^{1\gamma}$ is the probability of a single quantum transition to the ground state from the 2^3S_1 level, while $\mathscr{K}$ is the admixture coefficient of this level to the 2^1S_0 level due to the hyperfine interaction V_{HF}. These quantities are equal to:

$$w_{s3}^{1\gamma} = 1.0\cdot 10^{-3}\alpha^{11}Z_e^{10}m = 3\cdot 10^{-27}Z_e^{10}m,$$

$$\mathscr{K} = \langle 2\,{}^1S_0|V_{HF}|2\,{}^3S_1\rangle L_S^{-1} = \frac{7}{12}\left[\frac{I(I+1)}{2}\right]^{1/2}\frac{\alpha^4 Z^3 g m^2}{m_p L_S}$$

$$= 5.5\cdot 10^{-13}[I(I+1)]^{1/2}\frac{gmZ^3}{L_S},$$

where Z_e is the effective nuclear charge, Z is the nuclear charge, I is the nuclear spin, g is the gyromagnetic ratio, m_p is the proton mass, L_S is the energy difference, $L_S = \mathscr{E}(2^1S_0) - \mathscr{E}(2^3S_1)$.

All the formulas for the transition probabilities and for V_W contain the effective nuclear charge Z_e, while the formulas for the matrix elements V_{HF} contain the true charge Z. This is explained by the fact that the transition probabilities and V_W are determined by the external valence electrons, while the matrix elements are wholly determined by the behavior of the wave functions of the internal electrons near the nucleus. In the case of many-electron neutral atoms $Z_e \approx 1$, and in the case of multiply charged ions with a small number of electrons $Z_e \approx Z$.

The probability of transition, taking into account the circular polarization of the photons or the asymmetry with respect to the initial spin of the atom, is obtained from (4) by replacing $\mathbf{nh}$ by $\mathbf{ns}$ or $\mathbf{n\zeta}$ with the same value of P, where s and ζ are the photon and the atomic spins. The expression for P has the form[7]

$$P = P_1 + P_3, \quad P_l = \frac{2\gamma_l L_{Pl}^H}{(L_{Pl}^H)^2 + \Gamma_l^2/4 + \gamma^2}, \quad \gamma = \gamma_1 + \gamma_3, \quad \gamma_l = \left(\frac{w_{Pl}^{1\gamma}}{w_{S1}^{1\gamma}}\right)^{1/2}\delta_l,$$

$$\delta_l = \langle 2\,{}^1S_0|V_W|2\,{}^lP_1\rangle = \sigma_l\frac{bGm^3(\alpha Z_e)^4[2I(I+1)]^{1/2}}{32\pi}\left(\frac{R}{a_0}\right)^{-(\alpha Z_e)^2}$$

$$= \sigma_l bm\cdot 1.0\cdot 10^{-22}\left(\frac{R}{a_0}\right)^{-(\alpha Z_e)^2} Z_e^4[I(I+1)]^{1/2}, \quad \sigma_1 = 1, \quad \sigma_3 = 2^{1/2}, \tag{5}$$

$$L_{Pl}^H = \mathscr{E}(2\,{}^lP_1) - \mathscr{E}(2\,{}^1S_0), \quad \Gamma_1 = w_{P1}^{1\gamma} = 4\cdot 10^{-2}\alpha^5 Z_e^4 m$$

$$= 10^{-12}Z_e^4 m, \quad \Gamma_3 = w_{P3}^{1\gamma} = 10^{-1}\alpha^9 Z_e^8 m = 10^{-20}Z_e^8 m.$$

Here Γ_l is the width of the 2^lP_1 levels, $w_{Pl}^{1\gamma}$ are the probabilities of single quantum transitions from these levels to the ground state, R is the nuclear radius, a_0 is the Bohr radius[1)].

The difference between the energy levels L_{Pl}^H is evaluated taking into account the Zeeman shift of the sublevel of 2^lP_1 which has the value of the component of the total angular momentum $M = +1$ in the magnetic field H: $L_{Pl}^H = L_{Pl}^0 - \mu_0 H$. The sublevel with $M = 1$ is the sublevel which comes closer to the 2^1S_0 level as H increases. We note that we are considering such a situation when the Zeeman splitting is much greater than the hyperfine splitting: $\mu_0 H/\Delta_{HF} \gg 1$, and therefore the whole hyperfine multiplet behaves in a magnetic field as a single entity. For the same reason we neglect the Zeeman splitting of the hyperfine components of the 2^1S_0 level. In the situation which is of further interest to us we can also neglect the Zeeman splitting of the 2^3S_1 level, i.e., the dependence of L_S on H, since the distance L_S is sufficiently great ($\mu_0 H/L_S \ll 1$).

We note further that the magnetic field does not alter the transition probabilities in our case. Indeed, the transition $2^1S_1 \to 1^1S_0$ is allowed in the magnetic field only if the spin-orbit interaction is taken into account:

$$w_{S1}^{1\gamma}(H) \sim \mu_0 H(\alpha Z_e)^2 w_{S3}^{1\gamma}/L_S;$$

with the values of the intensity H of the magnetic field which we are considering we have $\mu_0 H(\alpha Z_e)^2/L_S \ll \mathscr{K}$. Finally we note that only the first term in (3) makes a contribution to P_1 (the singlet-singlet transition in (5)) while only the second term in (3) makes a contribution to P_3 (the singlet-triplet transition in (5))[2)]

The last term in (4) arises without the participation of the weak interaction under the influence of random electric fields, and under the experimental conditions must not exceed in magnitude the second term in (4). This imposes restrictions on the allowable value of the random fields E. The correlation $(\epsilon\mathbf{h})(\mathbf{nh})$ is a T-odd one and must vanish in the absence of the fields E and H, and therefore the magnitude of Q is proportional to H, E and to the level width Γ. In order of magnitude the coefficient Q is equal to[5]

$$Q = \sum_{l=1,2} Q_l, \quad Q_l \sim C_l\frac{(\mathbf{dE})(\mu_0\mathbf{H})\Gamma_l}{(L_{Pl}^H)^2 L_{Pl}^0}\left(\frac{w_{Pl}^{1\gamma}}{w_{S1}^{1\gamma}}\right)^{1/2}, \tag{6}$$

where $d = d_0/Z_e$, $d_0 = ea_0$, e is the electron charge, a_0 is the Bohr radius, μ_0 is the Bohr magneton.

For 2^1P_1 level the coefficient $C_1 \sim 1$, while for the 2^3P_1 level the coefficient $C_3 \sim (\alpha Z)^2$, since the mixing of the 2^3P_1 and 2^1S_0 levels occurs only with the participation of the spin-orbit interaction. Taking into account, moreover, that the transition $2^3P_1 \to 1^1S_0$ is forbidden and is lifted only by the spin-orbit interaction, we obtain $w_{P3}^{1\gamma} \sim (\alpha Z)^4 w_{P1}^{1\gamma}$ (cf., (3)) and $Q_3 \sim (\alpha Z)^4 Q_1$. This circumstance considerably relaxes the restrictions imposed on the external random electric fields when the levels are moved closer together by a magnetic field. The electric field also creates additional obstacles for observing the effect as a result of the Stark broadening of the 2^1S_0 level, and this leads to the single quantum decay with the emission of E1-photons.

If the M1- and E1-photons are not distinguished in the experiment this leads to an effective decrease in P. Therefore there arises a second condition imposed on E according to which the Stark broadening must not exceed the value of $w_{S1}^{1\gamma}$. This condition has the form[5]

$$\left(\frac{a_0 E}{Z_e L_{Pl}^H}\right)^2\frac{w_{Pl}^{1\gamma}}{w_{S1}^{1\gamma}} \ll 1 \tag{7}$$

and determines the allowable value of E when the experiment is performed without an external special magnetic field. The condition $Q \ll P$ then determines the allowable value of the random magnetic fields.

As a result of experiments with beams of ions passing through a foil (beam-foil spectroscopy) still another difficulty can arise: a correlation of the form $(\mathbf{Nn})(\mathbf{nh})$ appears where $\mathbf{N}$ is the normal to the foil (a pseudovector)[9]. The coefficient in front of this correlation may be of the order of unity, and therefore it can be eliminated without losses only for $P \sim 1$.

All the quantities in the expression for P can be conveniently rewritten in units of Γ_l:

$$P_l=\frac{2\beta_l x_l}{x_l^2+\beta^2+1},\quad x_l=\frac{2L_{Pl}^H}{\Gamma_l},\quad \beta=\beta_1+\beta_3,$$
$$\beta_l=\frac{2\gamma_l}{\Gamma_l}=\frac{2\delta_l}{(w_{S1}^{1\gamma}w_{Pl}^{1\gamma})^{1/2}} \tag{8}$$
$$\beta_1=7\cdot 10^8\frac{1}{Z^3Z_e^3}\frac{L_Sb}{mg}\left(\frac{R}{a_0}\right)^{-(\alpha Z_e)^2},$$
$$\beta_3=1.0\cdot 10^{13}\frac{1}{Z^3Z_e^5}\frac{L_Sb}{mg}\left(\frac{R}{a_0}\right)^{-(\alpha Z_e)^2}.$$

It is obvious that the 2^1S_0 and $2P_1$ can be moved closer together by a magnetic field by decreasing the value of L_{Pl}^H to $\Gamma_l(x_l \sim 1)$. If $\beta \gtrsim 1$, then the degree of asymmetry P can in this case be made to be of the order of unity. For $\beta \ll 1$ the value of $P_l^{max} = \beta_l$, as can be seen from (8). In the case of a hydrogenlike atom $\beta \sim 10^{-3}Z^{-3}$ [3], while in the case of two-electron atoms, as may be seen from (8), $\beta_l \gg 1$, if the charge Z_e is not too great, and for $\beta_l = x_l$ we obtain $P = 1$. The width of the range of values of H in which $P \sim 1$ is determined in accordance with (3) by the relation $\mu_0\Delta H = \gamma_l$, i.e.,

$$\Delta H/H=\gamma_l/\mu_0 H=\gamma_l/L_{Pl}^0.$$

Under such conditions the observation of single-quantum decay, i.e., the experimental statistics, becomes the principal problem. If one assumes that when the atoms are being excited the 2^1S_0 and 2^3P_1 levels are equally populated, then at the frequency of the single-quantum transition of interest to us intense radiation arises from the wing of the Breit-Wigner distribution for the 2^lP_1 level. This can considerably diminish the degree of asymmetry P. However, the metastable nature of the 2^1S_0 level provides a possibility of measuring its decay modes by the delay method, when after the excitation of the atoms a certain time elapses prior to the beginning of measurement during which the 2^lP_1 levels are vacated. The principal mode of decay of the 2^1S_0 level is the two-quantum decay with the probability $w^{2\gamma} = 2\cdot 10^{-20}Z_e^6 m$. The statistics of the experiment, i.e., the relative number of quanta of interest to us, is determined thus by the relationship

$$w_{S1}^{1\gamma}/w^{2\gamma}=0.5\cdot 10^{-31}(gm/L_S)^2Z_e^{10}(Z/Z_e)^6.$$

Here it is assumed that steps have been taken to prevent the two-quantum radiation from reaching the apparatus recording the asymmetry, i.e., that a filter is inserted which passes radiation in a relatively narrow frequency band $\Delta\omega$. In the presence of the filter the number of photons resulting from two-quantum decay which reach the recording apparatus is reduced by a factor of $(\Delta\omega/\omega_S)^2$ where ω_S is the frequency of the transition $2^1S_0 \to 1^1S_0$. The value of $\Delta\omega$ must be chosen so that this number would be smaller than the number of photons from single-quantum decay.

We now exhibit several specific examples.

1. <u>The He^3I atom.</u> For it $I = 1/2$, $g = 4.2$, $Z = 2$, $Z_e = 1.35$[3)], $L_S = 6.4\times 10^3\ cm^{-1} = 1.6\cdot 10^{-6}$ m, $L_{P1}^0 = 4.8\times 10^3\ cm^{-1} = 1.2\times 10^{-6}$ m, $L_{P3}^0 = 2.8\times 10^3\ cm^{-1} = 0.70\times 10^{-6}$ m, $w_{S1}^{1\gamma} = 0.8\times 10^{14}\ sec^{-1} = 1.1\times 10^{-35}$ m, $w_{S1}^{1\gamma}/w^{2\gamma}1.6\times 10^{-16}$, $\beta_1 = 0.2\times 10^3$ b, $\beta_3 = 1.0\times 10^6$ b.

In the absence of a magnetic field $P_1 = 0.2\times 10^{-3}$ b, $P_3 = 0.7\times 10^{-7}$ b. In the Weinberg model for the He^3 nucleus we obtain $b = b_W = 0.24$. Restrictions on the random fields are: $H_{max} < 10^7$ G, $E_{max} < 0.1$ V/cm. The degree of asymmetry P can be made to approach 1 in a magnetic field of $H \sim 10^7$ G. In this case the restriction on the random electric field becomes almost as strong as for hydrogen[5]: $E_{max} < 3\times 10^{-8}$ V/cm. Fields of intensity $H \sim 10^7$ G are at present unattainable. In the absence of a field the main obstacle to measurements is the smallness of the ratio $w_{S1}^{1\gamma}/w^{2\gamma}$, i.e., the small yield of the required photons.

2. <u>The C^{13}V ion[4)].</u> For it $I = 1/2$, $g = 1.4$, $Z = 6 \approx Z_e$, $L_S = 6.0\times 10^4\ cm^{-1} = 1.5\times 10^{-5}$ m, $L_{P1}^0 = 3.5\times 10^4\ cm^{-1} = 0.67\times 10^{-5}$ m, $L_{P3}^0 = 1.3\times 10^2\ cm^{-1} = 3.2\times 10^{-8}$ m, $w_{S1}^{1\gamma} = 2\times 10^{-8}\ sec^{-1} = 1.8\times 10^{-29}$ m, $w_{S1}^{1\gamma}/w^{2\gamma}2.5\times 10^{-13}$, $\beta_1 = 1.6$ b, $\beta_3 = 0.65\times 10^3$ b (in the Weinberg model $b = b_W = 0.24$ and $\beta_1 \sim 1$, $\beta_3 \gg 1$).

In the absence of a magnetic field $P_1 = 2.2\times 10^{-4}$ b, $P_3 = 1.8\times 10^{-4}$ b. The restrictions on the random fields are: $H_{max} < 6\times 10^5$ G, $E_{max} < 10^2$ V/cm.

In view of the anomalously small separation of the 2^1S_0 and 2^3P_1 levels (cf.,[11,12]) the carbon ion is the most favorable object for the measurement of the asymmetry P in a magnetic field. The Zeeman sublevel of the 2^3P_1 level with the value of the component of the angular momentum $M = +1$ crosses over with the 2^1S_0 level in the field of $H = 10^6$ G, and this is already on the threshold of being possible. In this case it is also necessary to take into account the fact that the 2^3P_0 and 2^3P_2 levels are also very close to 2^1S_0. The 2^3P_2 level also has a sublevel with the component $M = +1$, and therefore the sublevel which crosses with 2^1S_0 in the magnetic field of $H \sim 10^6$ G does indeed consist of a mixture of the sublevels of 2^3P_1 and 2^3P_2 with the same values of $M = +1$. For the determination of $L_P^{H_3}$ we must, thus, directly diagonalize the operator for the interaction with a magnetic field at the two aforementioned levels. The solution of the corresponding secular equation is

$$L_{P3}^H=L_{P3}^0-{}^1/_2\{(\Delta_P+3\mu_0 H)-(\Delta_P^2+\mu_0^2H^2)^{1/2}\},$$

where $\Delta_P \equiv \mathscr{E}(2^3P_2) - \mathscr{E}(2^3P_1)$. The width of the range of values of H for which $P \sim 1$, is equal to $\Delta H \approx 1$ G. Thus, there exists a quite narrow "resonance" in the asymmetry to hit which, apparently, is also an experimental problem. The restriction on the random electric fields now turns out to be the following: $E_{max} < 10^{-6}$ V/cm.

3. <u>The Cu^{65}XXVIII ion.</u> The 2^1S_0 and 2^3P_1 levels cross in addition to $Z = 6$ also at $Z = 29$ [11,12], and this corresponds to the two-electron copper ion. In this case $I = 3/2$, $g = 1.45$, $Z = 29 = Z_e$, $L_S = 0.75\times 10^{-4}$ m, $L_{P3}^0 = 1.0\times 10^{-7}$ m, $w_{S1}^{1\gamma} = 4.0\times 10^{-19}$ m $= 0.85\times 10^3$ sec^{-1}, $w_{S1}^{1\gamma}/w^{2\gamma} = 0.7\times 10^{-7}$, $\beta_1 = 2\times 10^{-2}$ b, $\beta_3 = 0.5\times 10^{-2}$ b.

Thus, in this case the application of an external magnetic field cannot yield the degree of asymmetry $P = 1$. In the absence of the field $P_1 = 1.4\times 10^{-4}$ b, $P_3 = 2.5\times 10^{-4}$ b. The limitations on the external field are: $H_{max} < 10$ G, $E_{max} < 10$ V/cm. The restriction on the magnetic field in this case is more rigorous than for the carbon ion. This is explained by the fact that the width of the 2^3P_1 level in going from $Z = 6$ to $Z = 29$ increases by a factor of 10^6 and the simulation of the nonconservation of parity in accordance with (6) is strengthened. However, it is possible to relax the limitation on H_{max} by strengthening the limitation on E_{max} by the same factor.

In conclusion we want to emphasize that the carbon ion $C^{13}V$ seems to us to be the most advantageous and to a sufficient degree unique object for measurements. Both the value of the magnetic field required to obtain complete asymmetry, and also the restriction on the magnitude of random electric fields do not appear to us to be unrealizable. We now make an estimate of the time required to acquire sufficient statistics in the experiment. The total number of recorded quanta is equal to $N = \nu\rho V\tau/\tau_0$, where ρ is the density of atoms in the required state, v is the working volume, τ is the time of observation, $\tau_0 = (w^{2\gamma})^{-1}$ is the time of irradiation, ν is a coefficient characterizing the efficiency of the experiment (the ratio of the number of recorded quanta to the number of those emitted). Assuming that for the observation of asymmetry it is necessary to have 10 "required" quanta, i.e., a total number of $\sim 10^{13}$ quanta, and assuming that $\rho = 10^{13}$ atoms/cm^3 (for larger values of ρ the lifetime of the $2\,^1S_0$ level is already determined by collisions), $V = 1$ cm^3, $\nu = 10^{-6}$ we obtain the observation time of $\tau \sim 1$ sec. Further we note that the frequency of the $2\,^1S_0 \to 1\,^1S_0$ transition in which we are interested lies in the x-ray region ($\omega \sim 330$ eV). Although this circumstance can make it more difficult to measure the degree of circular polarization of the radiation it ought not to be an obstacle in measuring the degree of asymmetry. And, finally, we make an estimate of the size of the transmission band $\Delta\omega$ of the filter required to reduce the number of recorded photons from the two-quantum decay. In the case of the $C^{13}V$ ion the required reduction is by a factor of at least 10^{12}. From this we obtain: $\Delta\omega/\omega_S \approx 10^{-6}$.

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1) The dependence on R appears when relativistic corrections are corrections are taken into account as a result of the singularity of the Dirac Coulomb functions in the neighborhood of the nucleus.

2) In principle the two terms in (3) can have different constants G_1 and G_2 if one takes into account the possibility of the appearance of a term of the form $\sigma_{\mu\nu}q_\nu$ in the electron current [5].

3) We have estimated the magnitude of Z_e for He^3I atoms by means of the formula $Z_e = [Z_e(2\,^1S_0)Z_e(2\,^3S_1)]^{1/2}$, where $Z_e(2\,^1S_0) = 1.21$, $Z_e(2\,^3S_1) = 1.51$ [10].

4) In the numerical data for the $C^{13}V$ and Cu^{65} XXVIII ions given below certain errors which occurred in [7] have been corrected.

1 Ya. B. Zel'dovich, Zh. Eksp. Teor. Fiz. 36, 964 (1959) [Sov. Phys.-JETP 9, 681 (1959)].
2 F. Curtis-Michel, Phys. Rev. 138B, 408 (1965).
3 A. N. Moskalev, ZhETF Pis. Red. 19, 229, 394 (1974) [JETP Lett. 19, 141, 216 (1974)].
4 M. A. Bouchiat and C. C. Bouchiat, Phys. Lett. 48B, 111 (1974).
5 Ya. I. Azimov, A. A. Ansel'm, A. N. Moskalev, and R. M. Ryndin, Zh. Eksp. Teor. Fiz. 67, 17 (1974) [Sov. Phys.-JETP 40, 8 (1975)].
6 I. B. Khriplovich, ZhETF Pis. Red. 20, 686 (1974) [JETP Lett. 20, 315 (1974)].
7 V. G. Gorshkov and L. N. Labzovskiĭ, ZhETF Pis. Red. 19, 768 (1974) [JETP Lett. 19, 394 (1974)] Preprint, Leningrad Nuclear Physics Institute, No. 120, 1974.
8 S. Weinberg, Phys. Rev. D5, 1412 (1972).
9 H. G. Berry, L. J. Curtis, D. G. Sills, and R. M. Schechtman, Phys. Rev. Lett. 32, 751 (1974).
10 H. Bethe and E. Salpeter, Quantum Mechanics of One- and Two-Electron Atoms, (Russ. Transl. Fizmatgiz, 1960).
11 G. L. Klimchitskaya and L. N. Labzovskiĭ, Opt. Spektr. 34, 633 (1973).
12 G. L. Klimchitskaya, Vestnik LGU (Bulletin of the Leningrad State University) 16, (1975).

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