culation of Z_{cr} for a large number of states of the discrete spectrum has appeared recently.[11]

[3)] The function $f(r/R)$, which cuts off the growth of $V(r)$ at $r<R$, is determined by the distribution of protons inside the nucleus. The following models of cutoff are used most frequently: I) $f(x)=1$ at $0<x<1$; II) $f(x)=(3-x^2)/2$, which corresponds to constant density of protons of form (3). According to the terminology used in Refs. 8, 9, and 11, these models are known as model I and model II, respectively.

[4)] The idea of the possibility of formation of a π-condensate in nuclei was expressed in Ref. 12. Detailed references to subsequent works and a discussion of the present status of the theory of the π-condensate can be found in Ref. 18.

[1]P. Gombas, Statistical theory of the atom and its application (Russian translation) IIL, 1951.
[2]N. H. March, Adv. Phys. 6, 1 (1957).
[3]D. A. Kirshnitz, Yu. E. Lozovik and G. V. Shpatakovskaya, Usp. Fiz. Nauk 117, 1 (1975) [Sov. Phys. Usp. 18, 649 (1975)].
[4]B. Müller and J. Rafelski, Phys. Rev. Lett. 34, 349 (1975).
[5]A. B. Migdal, V. S. Popov and D. N. Voskresenskiĭ, Pis'ma Zh. Eksp. Teor. Fiz. 24, 186 (1976); [JETP Lett. 24, 163 (1976)]; Zh. Eksp. Teor. Fiz. 72, 834 (1977) [Sov. Phys. JETP 45, 436 (1977)].
[6]I. Pomeranchuk and Ya. Smorodinsky, J. Phys. USSR. 9, 97 (1945).
[7]W. Pieper and W. Greiner, Z. Phys. 218, 327 (1969).
[8]V. S. Popov, Pis'ma Zh. Eksp. Teor. Fiz. 11, 254 (1970); [JETP Lett. 11, 162 (1970)]; Yad Fiz. 12, 429 (1970) [Sov. J. Nucl. Phys. 12, 235 (1971)].
[9]Ya. B. Zel'dovich and V. S. Popov, Usp. Fiz. Nauk 105, 403 (1971) [Sov. Phys. Usp. 14, 673 (1971)].
[10]A. I. Akhiezer and V. B. Berestetskiĭ, Kvantovaya élektrodinamika (Quantum Electrodynamics) Nauka, 1969.
[11]V. L. Eletskiĭ and V. S. Popov, Yad. Fiz. 25, 1107 (1977) [Sov. J. Nucl. Phys. 25, 587 (1977)].
[12]A. B. Migdal, Zh. Eksp. Teor. Fiz. 61, 2209 (1971) [Sov. Phys. JETP 34, 1184 (1972)].
[13]A. B. Migdal, Phys. Lett. 52B, 182 (1974).
[14]M. Gell-Mann and F. E. Low, Phys. Rev. 95, 1300 (1954).
[15]N. N. Bogolyubov and D. V. Shirkov, Vvedenie v teoriyu kvantovannykh polei (Introduction to the Theory of Quantized Fields) Nauka, 1976.
[16]J. Rafelski, V. Müller and W. Greiner, Nucl. Phys. B68, 585 (1974).
[17]M. Gyulassy, Phys. Rev. Lett. 33, 921 (1974); Nucl. Phys. A244, 497 (1975).
[18]A. B. Migdal, O. A. Markin, I. N. Mishustin and G. A. Sorokin, Zh. Eksp. Teor. Fiz. 72, 1247 (1977) [Sov. Phys. JETP 45, 654 (1977)].
[19]V. M. Galitskiĭ, Usp. Fiz. Nauk 120, 138 (1976) [Sov. Phys. Usp. 19, 769 (1976)]; Pis'ma Zh. Eksp. Teor. Fiz. 25, 175 (1977) [JETP Lett. 25, 160 (1977)].
[20]V. A. Karnaukhov and S. M. Polikanov, Pis'ma Zh. Eksp. Teor. Fiz. 25, 328 (1977) [JETP Lett. 25, 304 (1977)].
[21]M. S. Marinov and V. S. Popov, Zh. Eksp. Teor. Fiz. 67, 1250 (1974) [Sov. Phys. JETP 40, 621 (1975)] J. Phys. A8, 1575 (1975).
[22]V. S. Popov, V. L. Eletskiĭ and M. S. Marinov, Zh. Eksp. Teor. Fiz. 73, 1241 (1977) [Sov. Phys. JETP 46, 653 (1977)].
[23]L. N. Lipatov, Pis'ma Zh. Eksp. Teor. Fiz. 24, 179 (1976); 25, 116 (1977) [JETP Lett. 24, 157 (1976); 25, 104 (1977)]; Zh. Eksp. Teor. Fiz. 72, 411 (1977) [Sov. Phys. JETP 45, 216 (1977)].

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Parity nonconservation effects in two-photon transitions in hydrogen atoms

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Parity nonconservation effects in two-photon transitions $1s_{(1/2)}\to 2p_{(1/2)}$ and $2s_{(1/2)}\to 2p_{(3/2)}$ in a hydrogen atom arising if the neutral weak currents do not conserve parity are considered. The magnitude of the effects in the general case is 10^{-8}–10^{-9}. However, in the case of absorption of photons with equal energies and parallel or antiparallel momenta in the transition $1s_{(1/2)}\to 2p_{(1/2)}$ the magnitude of these effects increases and attains values of 10^{-4}–10^{-6}.

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The discovery in 1973 of weak neutral currents[1,2] opened up new possibilities for the study of weak interaction by means of looking for the effects brought about by this interaction of nonconservation of parity in atomic transitions. One can expect particularly large effects in transitions which in the absence of weak interactions are for some reason suppressed. There is available a large number of theoretical papers in which the effect of weak interaction on one-photon transitions in atoms is discussed (cf., the review articles[3,4]).

Also of interest are effects of parity nonconservation in two-photon atomic transitions. Experiments on induced two-photon transitions in an atom have received wide application recently due to progress in laser technology. The study of such transitions occurring as a result of a simultaneous absorption of two photons whose total energy is equal to the energy of the transition has a number of advantages compared to the study of the usual one-photon transitions. Among them are, for example, the extension of the range of energies of transitions

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which can be studied with the aid of existing lasers, the possibility of easily eliminating Doppler line broadening,[5] the high resolving power of lasers associated with this, etc.

In the present paper we study effects of parity nonconservation in two-photon transitions $1s_{1/2} \to 2p_{1/2}$ and $2s_{1/2} \to 2p_{3/2}$ in a hydrogen atom. The transition $1s_{1/2} \to 2p_{1/2}$ lies in the frequency range accessible to modern lasers (the two-photon transition $1s_{1/2} \to 2s_{1/2}$ which has the same energy as the transition $1s_{1/2} \to 2p_{1/2}$ has already been observed experimentally[6]). As regards the transition $2s_{1/2} \to 2p_{3/2}$, it lies within the range of centimeter radio waves convenient for investigation.

As a result of the weak interaction a mixing of the $2s_{1/2}$- and $2p_{1/2}$-states occurs. Therefore in the amplitude of the transition $1s_{1/2} \to 2p_{1/2}$ ($2s_{1/2} \to 2p_{3/2}$) there appears an admixture of the amplitude of the transition $1s_{1/2} \to 2s_{1/2}$ ($2p_{1/2} \to 2p_{3/2}$). The interference between the basic and the admixed amplitudes leads to the effects of parity nonconservation. In Sec. 1 we consider the transition $1s_{1/2} \to 2p_{1/2}$. The magnitude of the effects of parity nonconservation in this transition, generally speaking, is of the order of 10^{-8}. However, in the case of absorption of photons of equal frequencies and parallel or antiparallel momenta the effects of parity nonconservation are increased and attain values of the order of 10^{-6}–10^{-4}. In Sec. 2 we consider the transition $2s_{1/2} \to 2p_{3/2}$, in which the effects of parity nonconservation are of the order of 10^{-8}–10^{-9}. At the end of the paper we give a brief discussion of the influence of external fields on these transitions.

1. THE TRANSITION $1s_{1/2} \to 2p_{1/2}$

The weak interaction between electrons and the nucleus leads to parity nonconservation in atomic transitions due to the mixing of atomic levels with the same values of angular momenta but of different parity. In the case under consideration the basic role is played by the mixing by the weak interaction of the final $2p_{1/2}$ state with the $2s_{1/2}$ state since the distance between these levels (the Lamb shift) is exceptionally small. The Hamiltonian for the weak parity nonconserving interaction of the electron with the proton can be written in the form

$$H=-2^{-1/2}G\{\varkappa_1(\bar e\gamma_\mu\gamma_5 e)(\bar p\gamma_\mu p)+\varkappa_2(\bar e\gamma_\mu e)(\bar p\gamma_\mu\gamma_5 p)\} \tag{1}$$

and is characterized by two weak constants at present as yet unknown: $\varkappa_1$ and $\varkappa_2$ (we consider only the coupling of currents without derivatives).

The magnitude of the mixing of the $2p_{1/2}$- and the $2s_{1/2}$-levels δ_F depends on the total angular momentum of the atom F. According to calculations[1)] we have

$$\begin{aligned}\delta_0&=-\frac{3^{1/2}G}{32\cdot 2^{1/2}\pi}\frac{\alpha^4m^3}{\Delta_0}(\varkappa_1-3\varkappa_2)=-1.4\cdot10^{-11}(\varkappa_1-3\varkappa_2),\\ \delta_1&=-\frac{3^{1/2}G}{32\cdot 2^{1/2}\pi}\frac{\alpha^4m^3}{\Delta_1}(\varkappa_1+\varkappa_2)=-1.2\cdot10^{-11}(\varkappa_1+\varkappa_2),\end{aligned} \tag{2}$$

where $\Delta_F=E(2s_{1/2},F)-E(2p_{1/2},F)$ is the Lamb splitting of hyperfine levels with angular momentum F. In the case that the hyperfine interaction is neglected the mixing of the levels is characterized by the averaged coefficient δ which is given by the formula

$$\delta=\frac{1}{4}(\delta_0+3\delta_1)=-\frac{3^{1/2}G}{32\cdot2^{1/2}\pi}\frac{\alpha^4m^3}{\Delta}\varkappa_1=-1.2\cdot10^{-11}\varkappa_1 \tag{3}$$

and depends only on the single constant $\varkappa_1$.

If the phenomenon of mixing of levels is taken into account the amplitude A_p for the two-photon transition $1s_{1/2} \to 2p_{1/2}$ acquires an admixture of the amplitude A_s which characterizes the two-photon transition $1s_{1/2} \to 2s_{1/2}$. As a result the total amplitude A has the form

$$A=A_p-i\delta A_s. \tag{4}$$

We discuss the explicit form of the amplitudes A_p and A_s. The amplitude of the admixed two-photon transition $1s_{1/2} \to 2s_{1/2}$ has been calculated by many authors (cf., for example, Refs. 7, 8). We cite the expression for this amplitude which is obtained when it is calculated in the momentum representation by the method proposed by Gorshkov and Polikanov.[9] In the lowest order in the fine structure constant α which corresponds to the absorption of two electric dipole photons the amplitude A_s has the form

$$A_s=4\pi\alpha m^{-1}\mathbf{e}_1\mathbf{e}_2T_0, \tag{5}$$

where $\mathbf{e}_1$, $\mathbf{e}_2$ are the photon polarization vectors, m is the electron mass, T_0 is a scalar function the dependence of which on the frequencies of the absorbed photons ω_1, ω_2, ($\omega_1+\omega_2=3\alpha^2m/8$) is given by the formula

$$\begin{aligned}T_0&=2^6\cdot5\cdot2^{1/2}\left\{F(\omega_1)+F(\omega_2)+\frac{1}{5\cdot3^5}\right\},\\ F(\omega)&=\frac{1}{2}\frac{\xi^5(\xi^2-1)}{(\xi+1)^6(1+\xi/2)^6}\int_1^\infty\frac{t^{t+2}dt}{(t-x)^6},\\ x&=\frac{(1-\xi)(1-\xi/2)}{(1+\xi)(1+\xi/2)},\quad \xi=\left(1-\frac{2\omega}{\alpha^2m}\right)^{-1/2}.\end{aligned} \tag{6}$$

The graph of the function $T_0(\omega_1, \omega_2)$ is given in the diagram of Fig. 1. The lack of dependence of T_0 on the directions of the momenta of the photons $\mathbf{n}_1$ and $\mathbf{n}_2$ ($\mathbf{n}_i=\mathbf{k}_i/\omega_i$) is related to the dipole nature of the amplitude A_s.

The transition $1s_{1/2} \to 2p_{1/2}$ in the lowest order in α occurs due to the absorption of electric dipole and magnetic dipole ($E1\times M1$) or of the electric dipole and elec-

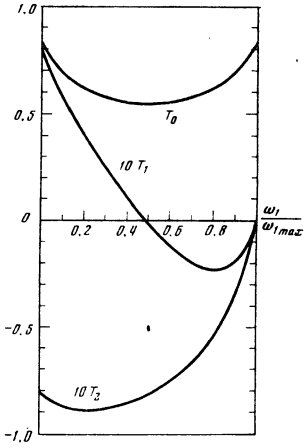

FIG. 1. Dependence of the amplitudes of two-photon absorption on the frequency of one of the photons.

tric quadrupole ($E1\times E2$) quanta. Therefore the amplitude for this transition A_p has an additional smallness $\sim\alpha$ compared with A_s. The dependence of A_p on the polarization vectors of the photons undergoing absorption $\mathbf{e}_1$, $\mathbf{e}_2$ and on the directions of their momenta $\mathbf{n}_1$, $\mathbf{n}_2$ can be written, on the basis of general requirements (symmetry, gradient invariance, etc.). This amplitude contains six independent structures and has the form

$$\begin{aligned}A_p=i4\pi\alpha^2m^{-1}\{&(\mathbf{e}_1\mathbf{e}_2)(\mathbf{n}_1\sigma)T_1(\omega_1,\omega_2)\\&+(\mathbf{e}_1\sigma)(\mathbf{e}_2\mathbf{n}_1)T_2(\omega_1,\omega_2)\\&+i([\mathbf{e}_1\times\mathbf{e}_2]\mathbf{n}_1)T_3(\omega_1,\omega_2)\\&+(1\leftrightarrow2)\}.\end{aligned}\tag{7}$$

Here σ are the Pauli matrices, the quantity $\frac{1}{2}\sigma$ plays the role of the electron angular momentum for the transition. The amplitude A_p is a pseudoscalar in contrast to A_s, and this is related to the different parity of the initial $1s_{1/2}$ and of the final $2p_{1/2}$ states.

We have calculated the functions T_i by means of the method proposed by Gorshkov and Polikanov.[9] As a result of this the following expressions are obtained for T_i:

$$\begin{aligned}T_1&=2^5\left(\frac{2}{3}\right)^{1/2}\left\{4F(\omega_1)+F(\omega_2)+\frac{1}{2\cdot3^5}-\frac{4}{3^5}\frac{\omega_1}{\alpha^2m}\right\},\\T_2&=2^5\left(\frac{2}{3}\right)^{1/2}\left\{4F(\omega_1)+F(\omega_2)-\frac{1}{3^5}\right\},\quad T_3=0,\end{aligned}\tag{8}$$

where $F(\omega)$ is defined in (6). The graphs of the functions T_0, T_1, and T_2 are given in the diagram of Fig. 1.

It can be seen from (8) that the amplitude of the $1s_{1/2}\to2p_{1/2}$ transition can be expressed in terms of the same functions of the frequency $F(\omega)$ as in the case of the amplitude for the $1s_{1/2}\to2s_{1/2}$ transition. Another important result is the equality $T_3=0$, i.e., the vanishing of that part of the amplitude A_p which does not contain the σ matrices. This means, in particular, that in the lowest approximation with respect to α two-photon transitions between hyperfine states $1s_{1/2}$ and $2p_{1/2}$ with total angular momentum $F=0$ are forbidden (terms containing σ give no contribution to these transitions). We explain why these transitions become forbidden.

The transition $1s_{1/2}(F=0)\to2p_{1/2}(F=0)$ occurs as a result of the absorption of $E1$- and $M1$-photons. Only three intermediate states with total angular momentum $F=1$ participate in it, specifically $1s_{1/2}$, $2p_{1/2}$ and $2p_{3/2}$, since the $M1$-transitions between states with different principal quantum numbers n are forbidden. It can be easily verified that the contributions of these intermediate states to the amplitude are in the ratio of $1:(-1/3):(-2/3)$, and therefore their total contribution is equal to zero. The transitions are no longer forbidden if one takes into account the difference in the energies of $2p_{1/2}$ and $2p_{3/2}$ or if one includes in the discussion $M1$-transitions between levels of different n or, finally, if one takes into account the absorption of $E2$- and $M2$-photons. In all these cases the amplitude acquires an additional smallness $\sim\alpha^2$, so that for making estimates one can assume

$$T_3\sim\alpha^2.\tag{9}$$

We now investigate the effects of parity nonconservation arising as a result of the interference of the A_p and A_s amplitudes. First we make several remarks of a general nature.

When the $1s_{1/2}\to2p_{1/2}$ transition is induced by laser radiation it is possible to separate the hyperfine components of the transition. The amplitudes of the individual hyperfine transitions are the matrix elements (5), (7) between the singlet χ_0 or the triplet χ functions. The admixed transition $1s_{1/2}\to2s_{1/2}$ occurs without a change of the total angular momentum of the atom F, therefore the parity violating correlations also appear only in transitions between the $1s_{1/2}$- and $2p_{1/2}$-levels with the same F. In this case the part of the amplitude A_p containing the matrices σ turns out to be responsible for the correlations associated with electron polarization. Conversely, the other part of A_p which does not depend on σ describes effects associated only with photon polarization. Since the principal transition $1s_{1/2}\to2p_{1/2}$ is suppressed by a factor of approximately α compared with the admixed transition $1s_{1/2}\to2s_{1/2}$ one can in the general case expect effects of parity nonconservation at the level of $\delta/\alpha\sim10^{-8}$–10^{-9}. By means of a special choice of experimental conditions the amplitude A_p can be suppressed thereby increasing the effects of parity nonconservation by several orders of magnitude. We consider two examples.

a) $\omega_1=\omega_2$, $\mathbf{n}_1=\mathbf{n}_2$. Such a configuration corresponds to the absorption of two photons from the same laser beam. In this case only one term remains in the amplitude A_p since $\mathbf{e}_1\cdot\mathbf{n}_2=\mathbf{e}_2\cdot\mathbf{n}_1=0$, and therefore the total amplitude taking into account the admixture of A_s has the form

$$A=i8\pi\alpha^2m^{-1}(\mathbf{e}_1\mathbf{e}_2)\,(\mathbf{n}_1\sigma)\,T_1-i\delta4\pi\alpha m^{-1}(\mathbf{e}_1\mathbf{e}_2)\,T_0.\tag{10}$$

The function $T_1(\omega_1,\omega_2)$, as can be seen from the diagram of Fig. 1, changes sign near the point $\omega_1=\omega_2$, and therefore at $\omega_1=\omega_2$ it is small:

$$T_1(\omega_1=\omega_2)=-0.94\cdot10^{-3},\quad T_0(\omega_1=\omega_2)=0.54.\tag{11}$$

Nonconservation of parity in such a transition leads, for example, to the atom becoming polarized in the final state when it was not polarized in the initial state, with the polarization vector being directed along the momentum of the absorbed photons. For the transition $F=1\to F=1$ the average value of the angular momentum of the atom in the final $2p_{1/2}$-state is given by the formula

$$\mathbf{P}=\langle\mathbf{F}\rangle=-\delta_1T_0\mathbf{n}_1/\alpha T_1=-0.94\cdot10^{-6}(\varkappa_1+\varkappa_2)\mathbf{n}_1.\tag{12}$$

The fluorescent photons emitted in the subsequent transition $2p_{1/2}\to1s_{1/2}+\gamma$ will also be polarized because of the polarization of the $2p_{1/2}$-state and their maximum circular polarization is given by formula (12).

b) $\omega_1=\omega_2$, $\mathbf{n}_1=-\mathbf{n}_2$. Such a configuration is utilized in the study of two-photon absorption without Doppler line broadening.[5] In this case only the term with T_3 remains in the amplitude A_p, while the total amplitude of the transition takes on the form

$$A=-8\pi\alpha^2m^{-1}([\mathbf{e}_1\mathbf{e}_2]\mathbf{n}_1)\,T_3-i\delta\cdot4\pi\alpha m^{-1}\mathbf{e}_1\mathbf{e}_2T_0.\tag{13}$$

Nonconservation of parity manifests itself in the dependence of the absorption cross section on the circular polarization of the photon (photons travelling towards each other must have the same signs of polarization). The relative difference of the absorption cross-section for the right- and left-polarized photons in the transition between $1s_{1/2}$- and $2p_{1/2}$-states with angular momentum F is given by the formula

$$\Delta \equiv \frac{\sigma(+)-\sigma(-)}{\sigma(+)+\sigma(-)} = -\delta_F \frac{T_0}{T_3 \alpha}. \tag{14}$$

Taking into account the high degree of suppression of the amplitude T_3 (cf., (9)) one can expect here effects at the level of 10^{-4}–10^{-5}. Naturally, the price paid for increasing the effects in the two aforementioned cases is the significant diminution in the probability of the transition.

The probability of the two-photon transition is given by

$$W = \frac{1}{4\omega_1\omega_2} J_1 J_2 A_F^2 \frac{1}{\Gamma}, \tag{15}$$

where J_i are the photon fluxes, Γ is the width of the final state. For given values of the density of atoms ρ, and of the working volume V one can, utilizing (15), estimate the value of the photon flux for which the time required for the observation remains realistic. The upper limit on the allowable values of the density ρ is determined by collisions which similarly to an external electric field lead to a mixing of $2s_{1/2}$- and $2p_{1/2}$-levels. In the case of configuration a) one obtains the restriction: $\rho \leq 10^{12}$ cm^{-3}. (A similar mixing in collisions with ions is more probable by a factor of α^2/v^2,[10] and the restriction on their density is: $\rho \leq 10^6$ ion/cm^3.) Setting the observation time to be of the order of one day, we obtain for the ideal case when the coefficient of efficiency of recording the absorption is equal to unity, we obtain

$$J \sim 2.5 \times 10^{21} \text{ cm}^{-2} \cdot \text{sec}^{-1}.$$

2. THE TRANSITION $2s_{1/2} \to 2p_{3/2}$

In this case the amplitude A_0 of the basic transition $2s_{1/2} \to 2p_{1/2}$ requires an admixture of the amplitude A_1 of the transition $2p_{1/2} \to 2p_{3/2}$ as a result of the mixing by the weak interaction of the initial state $2s_{1/2}$ with the state $2p_{1/2}$. Therefore the total amplitude is

$$A = A_0 + i\delta A_1. \tag{16}$$

The transition $2s_{1/2} \to 2p_{3/2}$ in the lowest order in α occurs by means of absorption of $E1$- and $M1$-quanta. (The amplitude, corresponding to the absorption of $E1$- and $E2$-quanta is smaller by a factor of $\omega r/\alpha \sim \alpha^2$.) In the sum over the intermediate states it is sufficient to take into account the contribution of states with the principal quantum number $n=2$: $2s_{1/2}$, $2p_{1/2}$, $2p_{3/2}$. Carrying out the summation over these states we have on neglecting the hyperfine interaction:

$$A_0 = K\Omega_{m'}^{+}\left\{\mathbf{a}_1(\sigma\mathbf{e}_2)\left(\frac{4}{\omega_1}+\frac{1}{\omega_2}\right)+\mathbf{a}_2(\sigma\mathbf{e}_1)\left(\frac{4}{\omega_2}+\frac{1}{\omega_1}\right) - \mathbf{e}_1\frac{3}{\omega_2}(\sigma\mathbf{a}_2)-\mathbf{e}_2\frac{3}{\omega_1}(\sigma\mathbf{a}_1)\right\}\varphi_m, \tag{17}$$

$$\mathbf{a}_i = [\mathbf{e}_i\mathbf{n}_i], \quad K = \frac{2\cdot 3^{1/2}\pi}{3m^2}\omega_1\omega_2,$$

where m and m' are the components of the angular momentum of the electron in the states $2p_{3/2}$ and $2s_{1/2}$, $\varphi_{m'}$ is the Pauli spinor

$$\Omega_m = \sum C^{s'_{1/2}, m}_{1,\mu;s_{1/2},\sigma}\varepsilon_\mu\varphi_\sigma;$$

ε_μ are the cyclic unit vectors.

The admixed transition $2p_{1/2} \to 2p_{3/2}$ occurs as a result of absorption of two $E1$-quanta. The principal contribution to the sum over the intermediate states is given by the state with the same principal quantum number and of opposite parity, i.e., $2s_{1/2}$. The amplitude of the mixed transition A_1 has the form

$$A_1 = b\Omega_{m'}^{+}\left\{\mathbf{e}_2\frac{(\sigma\mathbf{e}_1)}{\omega_1}+\mathbf{e}_1\frac{(\sigma\mathbf{e}_2)}{\omega_2}\right\}\varphi_{m'}, \quad b = \frac{12\pi\omega_1\omega_2}{\alpha m^2}. \tag{18}$$

We consider different hyperfine components of the transition. For this we project the wave functions of the bound electrons on states with the given total angular momentum of the atom F. The corresponding amplitudes are obtained from (17), (18) by the replacement

$$\begin{gathered}(\Omega_{m'}^{+}\mathbf{b}) \to (-1)^{M}\cdot 2^{-1/2}\{2b_{-M}\chi_0^{+} - i[\mathbf{b}\chi^{+}]_{-M}\} \quad (F=1),\\ (\Omega_{m'}^{+}\mathbf{b}) \to (-1)^{r}C^{2,M}_{1,r;1,\mu}b_{-r}\chi_\mu^{+}\ (F=2), \quad \varphi_{m'} \to \chi_0 \quad (F=0), \quad \varphi_{m'} \to \chi_{M'}\ (F=1),\end{gathered} \tag{19}$$

M' and M are the components of the total angular momentum of the atom in the initial and final states, χ_0 and χ are the wave functions for the total spin of the atom, $\mathbf{b}$ is an arbitrary vector.

Just as for the transition $1s_{1/2} \to 2p_{1/2}$ we consider two cases:

a) $\mathbf{n}_1 = \mathbf{n}_2 = \mathbf{n}$, $\omega_1 = \omega_2$. We calculate the polarization vector for atoms in the final state $\mathbf{P}_{F'F} = \mathrm{Sp}AA^{*}F/\mathrm{Sp}AA^{*}$, where F' and F are the angular momenta of the atom in the initial and final states. The vector $\mathbf{P}_{F'F}$ averaged over the polarizations of the quanta is directed along $\mathbf{n}$, with

$$\begin{gathered}\mathbf{P}_{F'F} = c_{F'F}\delta_F b\mathbf{n}/K,\\ \delta_0 b/K = -4.6\cdot 10^{-9}(\varkappa_1 - 3\varkappa_2), \quad \delta_1 b/K = -4.1\cdot 10^{-9}(\varkappa_1+\varkappa_2).\end{gathered} \tag{20}$$

Utilizing formulas (19) we obtain

$$c_{11} = {}^{11}/_{17}, \quad c_{12} = {}^{3}/_{7}, \quad c_{02} = -2. \tag{21}$$

b) $\mathbf{n}_1 = -\mathbf{n}_2$, $\omega_1 = \omega_2$, $\mathbf{e}_1 = \mathbf{e}_{R(L)}$, $\mathbf{e}_2 = -\mathbf{e}_1^{*}$. Evaluating

$$\Delta_{F'F} = \frac{\sigma(+)-\sigma(-)}{\sigma(+)+\sigma(-)},$$

where $\sigma(\pm)$ are the cross sections for the absorption of right- and left-polarized quanta, we have

$$\begin{gathered}\Delta_{F'F} = d_{F'F}\delta_F b/K,\\ d_{11} = 1, \quad d_{12} = 1, \quad d_{02} = -1,\end{gathered} \tag{22}$$

where $\delta_F b/K$ is determined by formula (20).

For the $0 \to 1$ transition A_1 vanishes at $\omega_1 = \omega_2$. Therefore in the $0 \to 1$ transition the effects of parity nonconser-

vation are not present in cases a) and b).

In conclusion, we discuss briefly the effect of external fields on the P-odd effects in the transitions $1s_{1/2} - 2p_{1/2}$ and $2s_{1/2} - 2p_{3/2}$. If a random electric field D is present in the experiment, then it also mixes states of opposite parity. The amplitude of such a transition A_D must not exceed the amplitude of the basic transition $A_p(A_0)$. If there is also present a random magnetic field H, then the interference between A_D and $A_p(A_0)$ leads to the appearance of effects simulating nonconservation of parity. For $H=0$ due to the interference between A_D and $A_p(A_0)$ a polarization vector $\mathbf{P}' \| \mathbf{n}\times\mathbf{D}$ appears but it is orthogonal to the polarization $\mathbf{P} \| \mathbf{n}$ brought about by the weak interaction. In the case of the transition $1s_{1/2} - 2p_{1/2}$ we have for the configuration a) the following restrictions: $D \leq D_0 = 10^{-2}$ V/m, $H \ll 3\times 10^{-2} D_0/D$ G. For the configuration b) in this transition $D \lesssim D_1 \sim 10^{-3}$ V/cm, $H \ll 3\times 10^{-2} D_1/D$ G. For the transition $2s_{1/2} - 2p_{3/2}$ the restrictions on the external fields have the form $D \leq D_2 = 1$ V/cm, $H \ll 2\times 10^{-5} D_2/D$ G.

In a field $H = 10^3$ G the levels $2s_{1/2}$ and $2p_{1/2}$ cross[11] and the magnitude of P-odd effect is increased by a factor of approximately 5 as a result of the increase in δ. However in this case it is apparently not possible to utilize this increase since there appears an induced electric field $\mathbf{D} = \mathbf{v}\times\mathbf{H}$. In the case of realistic values of $v \sim 10^{-5}$ $D \sim 2$ V/cm and $A_D \sim 10^3 \times A_p$, $A_D \sim 10\, A_0$.

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[1)]In this paper the system of units is utilized in which $\hbar = c = 1$.

[1]F. J. Hasert *et al.*, Phys. Lett. **B46**, 138 (1973).
[2]A. Benvenuti *et al.*, Phys. Rev. Lett. **32**, 800 (1974).
[3]V. A. Alekeseev, B. Ya. Zel'dovich, and I. I. Sobel'man, Usp. Fiz. Nauk **118**, 385 (1976) [Sov. Phys. Usp. **19**, 207 (1976)].
[4]A. N. Moskalev, R. M. Ryndin, and I. B. Khriplovich, Usp. Fiz. Nauk **118**, 409 (1976) [Sov. Phys. Usp. **19**, 220 (1976)].
[5]L. S. Vasilenko, V. N. Chebotaev, and A. V. Shishaev, Pis'ma Zh. Eksp. Teor. Fiz. **12**, 161 (1970) [JETP Lett. **12**, 113 (1970)].
[6]T. W. Hänsch, S. A. Lee, R. Wallenstein, and C. Wieman, Phys. Rev. Lett. **34**, 307 (1975).
[7]G. Breit and E. Teller, Astrophys. J. **91**, 215 (1940).
[8]B. A. Zon, N. L. Manakov, and L. P. Rappoport, Zh. Eksp. Teor. Fiz. **55**, 924 (1968) [Sov. Phys. JETP **28**, 480 (1969)].
[9]V. G. Gorshkov and V. S. Polikanov, Pis'ma Zh. Eksp. Teor. Fiz. **9**, 464 (1969) [JETP Lett. **9**, 279 (1969)].
[10]E. M. Purcell, Astrophys. J. **116**, 457 (1952).
[11]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 (1974)].

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Chemical radiative collisions

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A new type of elementary atomic collision event, namely, chemical radiative collisions, is examined. In each such event, the absorption (emission) of a photon is accompanied by a chemical reaction. A special case, namely, the substitution of an inert-gas atom R for a halogen atom in a molecule X_2 is analyzed. This process is accompanied by an optical transition of the quasimolecule RX_2 from a covalent to an ionic term. The associated light absorption coefficient is very large: it amounts to ~ 1 cm^{-1} for densities $N_{X_2} \sim N_R \sim 3\times 10^{19}$ cm^{-3}. In strong fields with amplitudes $E_0 \sim 10^7$ V/cm, the cross section for such chemical radiative collisions may be up to ~ 100 Å^2 and, at still higher light intensities, it ceases to depend on E_0. The results of the first experiments in this field are reported. Measurements of the absorption coefficient and of the corresponding wavelength range are found to be in agreement with theoretical predictions. Possible applications of the effect are discussed.

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1. We have investigated chemical reactions between molecules during collisions in a light field. In these previously unexplored reactions, which we shall refer to as "chemical radiative collisions" (or, briefly, chemical RC), the field participates not in the preparation of the states of the colliding particles X and Y, but in the elementary event itself. In general, a radiative collision (RC) is commonly understood to mean an elementary event in which a photon is emitted (or absorbed) during the collision, so that the conservation of energy can be satisfied during an inelastic transition of the XY molecule. By chemical RC we shall understand events in which the result of the application of the light field is that the collision produces not only a change in the structure of the electron shells of the colliding particles[1–6] but, in addition, the atoms form new bound states, i.e., a chemical reaction occurs during the collision. The simplest example of chemical RC are the well-known photoassociation and photodissociation reactions:

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