Polarization of nuclei in resonance ionization of atoms

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Resonance ionization of atoms via hfa states leading to polarization of the ion nuclei is considered for the case in which circularly polarized radiation is used to excite or ionize the atoms. The requirements on the frequency, monochromaticity, polarization, and strength of the light field and on the spin hfs splitting of the atomic levels are elucidated. Some general formulas for calculating the polarization of the ion nuclei as a function of the Stokes parameters of the radiation are obtained. The results of numerical calculations of the polarization of the nuclei of ions from resonance ionization of hydrogen and lithium atoms are presented.

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1. INTRODUCTION

It is well known that polarized electrons may be produced by multiphoton resonance ionization of atoms.[1] We recently called attention to the possibility that polarized nuclei may also be produced in that process.[2] The nuclei of the ions produced by resonance ionization will always be polarized if at least one of the resonant states is a spin hfs state and the atom is either excited or is ionized from the excited state by circularly polarized light. It would seem possible to develop a very widely usable method for obtaining beams of polarized nuclei using the resonance ionization of atoms by laser light. The problem of finding new methods for polarizing nuclei is a timely and urgent one. The methods used now, which are based on the spatial separation in a nonuniform magnetic field of atoms in different spin hfs states actually allows one only to obtain beams of polarized protons or helium nuclei.[3] We note that there have been recently discussed not only the resonance ionization of atoms,[2] but also two other methods for obtaining polarized nuclei using laser radiation: the vaporization of a previously polarized target[4] and the selective quenching of metastable states of the hydrogen atom.[5]

In this paper we discuss the basic conditions for obtaining ions with polarized nuclei by resonance ionization of atoms.

2. FUNDAMENTAL CONDITIONS FOR NUCLEAR POLARIZATION

Multiphoton resonance ionization is a process in which an atomic electron undergoes a transition from the ground state to the continuum, while at the same time two or more photons ($K_0 \geqslant 2$) are absorbed under such conditions that the energy of some smaller number of photons ($K_1 < K_0$) is close to the transition energy to some excited bound state. The atom may be excited and the excited atom ionized by different radiation fields having different frequencies, polarizations, and strengths. As will be shown below, the use of different fields makes it possible to maximize the nuclear polarization.

Since the operator $d\mathscr{E}$ for the interaction of the atom with the external field (d is the electron dipole moment) is independent of the nuclear variables, it is necessary that either the ground state or the excited state (or both) be spin hfs states if the nucleus is to be polarized. Another necessary condition is that either the exciting radiation or the ionizing radiation be circularly polarized. When an unpolarized atom absorbs one or more circularly polarized photons, the projection of the angular momentum of the atom or a given direction will always be nonuniformly distributed, i.e., the atom will be polarized to a certain extent. Thus, resonance absorption by atoms gives rise to polarized electrons and to ions with polarized nuclei.

Let ΔE and $\Delta E'$ be energies characteristic of the hfs splittings of the ground and resonant states. The condition for selective transitions between the hfs sublevels can be written in the form

$$\Delta \equiv \max(\Delta E, \Delta E') > \gamma \equiv \max(\Gamma, \Delta\omega). \qquad (1)$$

Here $\Delta\omega$ is the effective spectral width of the radiation, which depends on the spectral properties of the radiation employed and on the multiplicity K_1 of the resonance, and Γ is the reduced width of the excited and ground levels, for which various physical causes may be responsible. The principal broadening mechanisms are spontaneous decay of the excited levels, field and ionization broadening in the strong external field, inhomogeneous and Doppler broadening, and the collision of atoms with one another. Collisions can practically always be neglected in experiments with atomic beams. The Doppler broadening may also be comparatively small if the atomic and light beams are perpendicular to one another. In addition to the natural width, therefore, we need take into account only effects associated with the perturbation of the atomic levels by the external field. Except for some obvious estimates that we shall present in the Conclusion, however, we shall not be concerned in this paper with the effects of broadening and shifting of the levels in the external field. We shall assume that the perturbation of the resonant states is smaller than or of the order of the natural widths γ of the states. This permits us, in particular, to treat the transitions of the atom induced by the external fields in the lowest nonvanishing order (K_0) of perturbation theory. This approach is suitable from the practical point of view: as will be seen below, if the perturbation exceeds the natural width, the polarization will be smaller than in the case examined here.

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The resonance ionization process in which hfs states are selectively excited by circularly polarized light has been recently studied experimentally for the cesium atom.[6,7] In those studies the measured quantity was the polarization of the electrons, the polarization of the nuclei not being investigated. Nevertheless certain conclusions concerning the polarization of the nuclei can be drawn from the electron polarization observed in these experiments. Thus, it was shown experimentally[6,7] that the polarization of the electrons is higher for transitions between fine structure states than for transitions between the corresponding hfs states. This result is obviously associated with the fact that in the hfs scheme the nucleus acquires a definite fraction of the angular momentum of the absorbed photons. The decrease in the polarization of the electrons is consequently to be attributed to polarization of the nuclei.

Let us illustrate the possibility of polarizing nuclei within the framework of the formulated conditions, using the hydrogen atom as the simplest example.[8] The transition scheme between the hfs sublevels of the ground state ($1S_{1/2}$), those of the excited state ($2P_{1/2}$), and the continuum states for the case of a right-hand polarized exciting field $\mathscr{E}$ is shown in Fig. 1. If the frequency ω of the exciting field is close to the transition energy from the $1S_{1/2}$ ($F=0$) level to the $2P_{1/2}$ ($F=1$) level, the only possible intermediate state will be the $2P_{1/2}$ ($F=1$, $m_F=1$) state, where m_F is the projection of the total angular momentum F of the atom on the z axis. The wave function of this state is the following product of electronic and nuclear wave functions:

$$|I={}^1/_2,\ m_I={}^1/_2\rangle\,|j={}^1/_2,\ m_j={}^1/_2\rangle, \tag{2}$$

where m_I and m_j are the projections onto the z axis of the nuclear spin I and the electron angular momentum j. The proton is 100% polarized in this state, and therefore also in the continuum. We note that the transition between the $1S_{1/2}$ ($F=0$) and $2P_{1/2}$ ($F=0$) states is not energy wise forbidden, but is forbidden by the $\Delta m_F=1$ selection rule. Hence the nuclear polarization reaches 100% despite the fact that the hfs of the $2P_{1/2}$ level is not resolved, the radiative widths of the hfs sublevels exceeding the hyperfine splitting.[9]

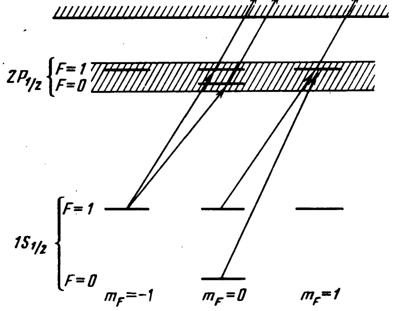

FIG. 1. Level scheme showing transitions between hfs sublevels of the $1S_{1/2}$ and $2P_{1/2}$ states under the action of two right-hand polarized fields. The hatching indicates the edge of the continuum and the width of the excited level.

3. THE DENSITY MATRIX AND THE DEGREE OF NUCLEAR POLARIZATION

Now let us derive general formulas for the nuclear polarization incident to resonance ionization. For simplicity we shall consider only atoms having just one electron above the filled shells.

Suppose that an atom initially in a state with the quantum numbers L, J, F, and m_F (where L is the electron orbital angular momentum, J is the total electron angular momentum, F is the total angular momentum of the atom, and m_F is the projection of the total angular momentum of the atom onto the z axis) absorbs a photon of frequency ω and undergoes a transition to a resonant state with the corresponding quantum numbers L', J', f, and m_f and then absorbs a second photon of frequency ω' and becomes ionized.

The expression for the transition amplitude for this process in the second order of perturbation theory is

$$A_\mu(\omega; F, m_F; L''M\sigma) = \sum_{f,m_f} (E_F+\omega-E_f)^{-1} \times\langle L''M\sigma I\mu|d\mathscr{E}'|L'J'fm_f\rangle\langle L'J'fm_f|d\mathscr{E}|LJFm_F\rangle. \tag{3}$$

Here $\mathscr{E}$ and $\mathscr{E}'$ are the strengths of the exciting and ionizing fields, L'' is the orbital angular momentum of the electron in the continuum, and M, σ, and μ are the projections of the orbital angular momentum, the electron spin, and the nulcear spin, respectively. We shall take the energies E_F and E_f as complex in order to take account of the natural widths of the atomic levels. We recall that the perturbations of the resonance states by the laser radiation are assumed to be smaller than the natural widths of those states.

As is well known,[9] the largest of the matrix elements for the transition from the resonance state to the continuous spectrum are those for which the orbital angular momentum L'' is larger than L': $L''=L'+1$. Taking only these matrix elements into account, we can write the following expression for the polarization density matrix of the nucleus:

$$\rho_{\mu\mu'}(\omega) = \sum_{FF',m_Fm_F',M\sigma} A_\mu(\omega; Fm_F; L'+1, M\sigma)\,\rho^{(0)}_{Fm_F,F'm_F'} \times A_{\mu'}(\omega; F'm_F'; L'+1, M\sigma), \tag{4}$$

where $\rho^{(0)}$ is the density matrix for the atom in the initial state. The matrix (4) is defined here to within a factor, which is fixed by the normalization condition $\mathrm{Tr}\rho=1$.

If we assume that the initial states are populated uniformly (with equal probability) and that the off-diagonal elements of $\rho^{(0)}$ have random phases, we can use standard methods of the theory of angular momenta[10] to obtain the following expression for $\rho_{\mu\mu'}$:

$$\rho_{\mu\mu'}(\omega) = \frac{1}{2I+1}\sum_{k\varkappa} C^{I\mu'}_{I\mu k\varkappa}\frac{a_{k\varkappa}(\omega,\mathscr{E},\mathscr{E}')}{a_{00}(\omega,\mathscr{E},\mathscr{E}')}, \tag{5}$$

$$a_{k\varkappa}(\omega,\mathscr{E},\mathscr{E}') = \sum_{Fff'jj'\lambda\lambda'} (-1)^{f-F+j}[(E_F+\omega-E_f)(E_F^*+\omega-E_{f'}^*)]^{-1}$$

$$\times(2F+1)(2f+1)(2f'+1)[(2j+1)(2j'+1)(2k+1)]^{1/2}C^{k\varkappa}_{j\lambda j'\lambda'}P_{j\lambda}(\mathscr{E})P_{j'\lambda'}(\mathscr{E}')$$

$$\times\begin{Bmatrix} j & 1 & 1 \\ F & f' & f \end{Bmatrix}\begin{Bmatrix} 1 & J & J \\ I & f & F \end{Bmatrix}\begin{Bmatrix} 1 & J & J' \\ I & f' & F \end{Bmatrix}\begin{Bmatrix} L' & 1 & L'+1 \\ 1 & L' & j' \end{Bmatrix}$$

$$\times\begin{Bmatrix} j' & L' & L' \\ {}^1/_2 & J' & J' \end{Bmatrix}\begin{Bmatrix} J' & I & f \\ j' & k & j \\ J' & I & f' \end{Bmatrix}. \tag{6}$$

Here the

$$P_{j\lambda}(\mathscr{E})=\frac{1}{\mathscr{E}^2}\sum_{pq}C_{1q1p}^{j\lambda}\mathscr{E}_p\mathscr{E}_q^* \tag{7}$$

are the polarization tensors for the exciting and ionizing radiations, and the summation indices in (5) and (6) assume, as usual, values that satisfy the triangle rule.[10] The components of $P_{j\lambda}$ that do not vanish in the natural coordinate system (the coordinate system with the z axis in the direction of the wave vector and the x axis in the direction of maximum linear polarization) are

$$P_{00}^{(c)}=-1/\sqrt{3},\quad P_{10}^{(c)}=A/\sqrt{2},\quad P_{20}^{(c)}=-1/\sqrt{6},\quad P_{2,\pm2}^{(c)}=l/2, \tag{8}$$

where A and l are the degrees of circular and linear polarization, respectively, and are given in terms of the Stokes parameters as follows[11]: $A=\xi_2$, and $l=(\xi_1^2+\xi_3^2)^{1/2}$. In an arbitrary coordinate system we have

$$P_{j\lambda}=\sum_{\nu}P_{j\nu}^{(c)}D_{\nu\lambda}^{j}(\alpha,\beta,\gamma),$$

where α, β, and γ are the Euler angles for the transformation from the natural coordinate system to the coordinate system under discussion and $D_{\nu\lambda}^{j}$ is the matrix effecting that transformation.

It is evident from the formulas given above that the matrix (5) will be diagonal if the exciting and ionizing radiations propagate along the same line and both are circularly polarized. In this case we take the quantization axis in the propagation direction of the waves. In all other cases the matrix (5) is nondiagonal.

In the general case, when we allow for the fact that the exciting and ionizing fields are not strictly monochromatic, formulas (5) and (6) for the density matrix ρ become much more complicated. If we assume for simplicity, however, that at least one of the fields (which are not correlated with one another) is a radiation field with random phase uniformly distributed on the interval $[-\pi,\pi]$ we can transform (5) to the form

$$\begin{gathered}\rho_{\mu\mu'}=\frac{1}{2I+1}\sum_{k\varkappa}C_{I\mu k\varkappa}^{I\mu'}\frac{b_{k\varkappa}(\mathscr{E},\mathscr{E}')}{b_{00}(\mathscr{E},\mathscr{E}')},\\ b_{k\varkappa}(\mathscr{E},\mathscr{E}')=\int d\omega g(\omega)a_{k\varkappa}(\omega,\mathscr{E},\mathscr{E}'),\end{gathered} \tag{9}$$

where $g(\omega)$ is the effective shape of the exciting spectrum line.

The density matrix $\rho_{\mu\mu'}$ completely determines the polarization of the ensemble of nuclei. In speaking of the polarization of the nuclei one frequently has in mind the average value of the projection of the nuclear spin onto some fixed direction:

$$P=\frac{1}{I}\sum_{\mu}\mu\rho_{\mu\mu}. \tag{10}$$

On substituting this definition into (8) we easily obtain the following expression for the polarization:

$$P=\frac{1}{3}\left(\frac{I+1}{I}\right)^{1/2}\frac{b_{10}(\mathscr{E},\mathscr{E}')}{b_{00}(\mathscr{E},\mathscr{E}')}. \tag{11}$$

Let us consider a few special cases in which the formulas for the polarization are somewhat simpler.

1. The effective width γ is small as compared with the hfs splitting for both the ground and the resonance states:

$$\gamma\ll\Delta E,\ \Delta E'. \tag{12}$$

This means that by suitably choosing the frequency ω we can bring about a selective transition between specified hfs components: F_0-f_0. In this case the transitions between all the other hfs components can be neglected because they are not in resonance. Such conditions can be brought about, for example, for the resonance lines of alkali atoms.[6,7,12] In this case, after cancelling common factors, we obtain the following equation from Eq. (6):

$$\begin{gathered}b_{k\varkappa}(\mathscr{E},\mathscr{E}')=\sum_{j\lambda j'\lambda'}(-1)^{j}[(2j+1)(2j'+1)(2k+1)]^{1/2}C_{j\lambda j'\lambda'}^{k\varkappa}P_{j\lambda}(\mathscr{E})P_{j'\lambda'}(\mathscr{E}')\\ \times\begin{Bmatrix} j & 1 & 1\\ F_0 & f_0 & f_0\end{Bmatrix}\begin{Bmatrix} L' & 1 & L'+1\\ 1 & L' & j\end{Bmatrix}\begin{Bmatrix} j & 1 & 1\\ 1/2 & J' & J'\end{Bmatrix}\begin{Bmatrix} J' & I & f_0\\ j' & k & j\\ J' & I & f_0\end{Bmatrix}.\end{gathered} \tag{13}$$

2. The hfs splitting of the resonant state is much smaller than the width γ, which, in turn, is considerably smaller than the hfs splitting of the ground state:

$$\Delta E\gg\gamma\gg\Delta E'. \tag{14}$$

In this case, by properly choosing the frequency ω we can ionize the atom to a state which has a definite value F_0 of the total angular momentum and is in resonance with all the hfs components of the excited level. This situation obtains, for example, for the hydrogen atom (see Section 2 above).

Assuming that the relation $E_j\approx E_{J'}$ holds for all the f values, we obtain the following equation from Eq. (5):

$$\begin{gathered}b_{k\varkappa}(\mathscr{E},\mathscr{E}')=\sum_{j\lambda j'\lambda'}(-1)^{k}[(2j+1)(2j'+1)(2k+1)]^{1/2}\\ \times C_{j\lambda j'\lambda'}^{k\varkappa}P_{j\lambda}(\mathscr{E})P_{j'\lambda'}(\mathscr{E}')\begin{Bmatrix} L' & 1 & L\\ 1 & L' & j'\end{Bmatrix}\\ \times\begin{Bmatrix} j' & L' & L'\\ 1/2 & J' & J'\end{Bmatrix}\begin{Bmatrix} I & I & k\\ J & J & F_0\end{Bmatrix}\begin{Bmatrix} J & J' & 1\\ J & J' & 1\\ k & j' & j\end{Bmatrix}.\end{gathered} \tag{15}$$

3. If the condition opposite to (12) obtains so that the natural width exceeds the hfs splitting, i.e., if

$$\gamma\gg\Delta E,\ \Delta E', \tag{16}$$

then from Eqs. (5) and (8) we find that $b_{k\varkappa}\sim\delta_{k0}\delta_{\varkappa0}$ and

$$\rho_{\mu\mu'}=\delta_{\mu\mu'}/(2I+1). \tag{17}$$

In view of the general considerations presented above in Section 1, it is obvious that in this case the nuclei will not be polarized at all.

4. DEPENDENCE OF THE NUCLEAR POLARIZATION ON THE POLARIZATIONS OF THE EXCITING AND IONIZING FIELDS

Let us examine the polarization of the nuclei in more detail for the case in which the wave vectors of the exciting and ionizing fields have the same direction (parallel to the z axis) but the two fields are arbitrarily polarized.

If we assume that the exciting and ionizing fields are not correlated with one another we can express the polarization P in general form in terms of the Stokes parameters ξ_α ($\alpha=1,2,3$) for the ionizing field ($\xi_{\alpha i}$)

and the exciting field, $(\xi_{\alpha e})$.[11] As is well known, the Stokes parameter ξ_2 specifies the degree of circular polarization of the radiation, while the parameters ξ_3 and ξ_1 specify the linear polarization along the x axis and along an axis in the xy plane making an angle of $\pi/4$ with the x axis; moreover, we have $\sum_\alpha \xi_\alpha^2 = 1$.

It follows from the general formulas of the preceding section that under the conditions specified above the polarization of the nuclei can be expressed in the form

$$P=\frac{m_{++}(\xi_{2i}+\xi_{2e})+m_{-+}(\xi_{2e}-\xi_{2i})}{N_{++}(1+\xi_{2i}\xi_{2e})+N_{-+}(1-\xi_{2i}\xi_{2e})+[2N_{00}-N_{-+}-N_{++}](\xi_{1e}\xi_{1i}+\xi_{3e}\xi_{3i})}\cdot \tag{18}$$

The constants occuring in this expression have visualizable meanings: m is the average value of the projection of the nuclear spin, N is the total number of ions, the subscripts $++$ and $-+$ refer to the cases in which both fields are right-hand polarized $(++)$ and in which the ionizing field is left-hand polarized and the exciting field, right-hand $(-+)$, and finally, N_{00} is the number of ions produced in two-photon resonance ionization of the atom by two fields linearly polarized in the same direction.

Formula (18) clearly shows that the polarization of the nuclei can only be due to the circular component of either the exciting or the ionizing field, since $P=0$ if $\xi_{2e}=\xi_{2i}=0$. The ionization of atoms by two linearly polarized fields does not lead to polarization of the nuclei. We note that this last conclusion does not depend on the assumption, used in the present section, that the wave vectors of the exciting and ionizing radiation are parallel, since in the dipole approximation the physical quantities do not depend on the propagation direction of the linearly polarized radiation (for fixed direction of the polarization vector). The linear polarizations of the exciting and ionizing fields, which are determined by the parameters ξ_{1e} and ξ_{3e}, and ξ_{1i} and ξ_{3i}, respectively, affect only the total number of nuclei produced (and thereby the numerical value of P).

On the other hand, Eq. (18) illustrates the fact that the exciting and ionizing fields are to a considerable extent equally important as regards their polarization properties. In particular, the nuclei will be polarized $(P\neq 0)$ even if only one of the two parameters ξ_{2e} and ξ_{2i} differs from zero. Thus, to obtain polarized nuclei it is sufficient that one of these fields be circularly polarized. This is obviously correct when condition (1) for selective excitation of the hfs states is satisfied. If the hfs of the excited state is not resolved [condition (14)], the equation giving the polarization P in terms of the Stokes parameters retains the general form of Eq. (18).

Further simplification of the general formula (18) is possible in the case of the hydrogen atom, in which the nuclear spin and the electron angular momentum in the ground state are both equal to $\frac{1}{2}$. In this case the polarization can be factored, being expressible as a product of two functions f_1 and f_2 of which one depends only on a quantum number (F) characterizing the initial state, and the other, only on a quantum number (j') characterizing the resonant state:

$$P=f_1(F)f_2(j'). \tag{19}$$

The factor $f_1(F)$ has the form

$$f_1(F)=2\Big[\sum_{m_I=\pm 1/2} m_I\big(C^{F,m_I-1/2}_{1/2,-1/2;1/2m_I}\big)^2\Big]\times\Big[\sum_{m_I=\pm 1/2}\big(C^{F,m_I-1/2}_{1/2,-1/2;1/2m_I}\big)^2\Big]^{-1}=\begin{cases}1, & F=0\\ -1/3, & F=1\end{cases}. \tag{20}$$

This means that for the hydrogen atom the polarization of the protons resulting from transitions from the $F=1$ initial state will (other things being equal) be one-third of the polarization of protons produced by ionization from the $F=0$ state and will have the opposite sign, regardless of the polarization characteristics of the exciting and ionizing fields.

The factor $f_2(j')$ in Eq. (19) reflects the dependence of the proton polarization on the choice of the resonant state. It also includes the dependence on the Stokes parameter ξ. For $j'=\frac{1}{2}$ we have

$$f_2(1/2)=\frac{\xi_{2e}+\xi_{2i}+\lambda(\xi_{2e}-\xi_{2i})}{1+\xi_{2e}\xi_{2i}+\lambda(1-\xi_{2e}\xi_{2i})}, \tag{21}$$

where λ is a certain constant for which a definite expression in terms of the transition matrix elements can be given. Its numerical value is $\lambda\approx 0.65$. An interesting fact to be seen from Eq. (21) is that in the case of resonance with the $P_{1/2}$ level the proton polarization is entirely independent of the linear polarization of the light. In the special case of (right-hand) circular polarization of the exciting radiation $(\xi_{2e}=1)$ we have $f_2(\frac{1}{2})=1$.

The expression for the factor $f_2(\frac{3}{2})$ for the case of resonance with the $P_{3/2}$ level is more complicated than that for $f_2(\frac{1}{2})$:

$$f_2(3/2)=\frac{A_1(\xi_{2e}+\xi_{2i})+A_2(\xi_{2e}-\xi_{2i})}{B_1(1+\xi_{2e}\xi_{2i})+B_2(1-\xi_{2e}\xi_{2i})+B_3(\xi_{1e}\xi_{1i}+\xi_{3e}\xi_{3i})}; \tag{22}$$

here A_1, A_2, B_1, B_2, and B_3 are certain constants. We shall not give the expression for these constants in terms of the matrix elements, but shall give only their numerical values: $A_1\approx 0.41481$, $A_2\approx 0.01975$, $B_1\approx 0.65185$, $B_2\approx 0.15802$, and $B_3\approx 0.18963$.

For the special cases $\xi_{2e}=\xi_{2i}=1$ and $\xi_{2e}=-\xi_{2i}=1$ we have

$$[f_2(3/2)]_{++}=-7/11,\qquad [f_2(3/2)]_{-+}=-1/8,$$

respectively.

Figure 2 shows plots of the proton polarization vs one of the Stokes parameters for fixed values of the others.

No particular difficulties are encountered in obtaining analogous relationships for other atoms. Below we give data for the lithium atom as an example. If should be noted that from the experimental point of view the alkali atoms provide more suitable targets for obtaining polarized nuclei than does the hydrogen atom. The reason is that for hydrogen the transition to the first excited state lies in the ultraviolet, while for alkali atoms it lies in the visible, and it is well known that it is much easier to obtain the intense laser light necessary for exciting the atoms in the visible than in the ultraviolet. As was noted above, the alkali atoms satisfy the first condition for achieving selective transitions between the sublevels of the ground and resonant states; hence

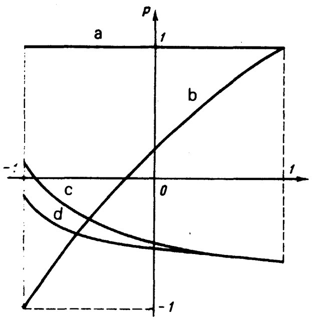

FIG. 2. Polarization P of the protons from resonance ionization of the hydrogen atom from the $1S_{1/2}$ ground state ($F=0$) via the $2P_{1/2}$ (curves a and b) and $2P_{3/2}$ (curves c and d) excited levels. For curves a and d, $\xi_2=\xi_{2i}$ and $\xi_{2e}=1$; for curves b and c, $\xi_2=\xi_{2e}$ and $\xi_{2i}=1$.

formula (18) is to be used for the calculations. The function $P(\xi_{2e})$ for the fixed value $\xi_{2i}=1$ and the function $P(\xi_{2i})$ for $\xi_{2e}=1$ are linear fractional functions, as in the case of the hydrogen atom. Figure 3 shows these functions for resonance transitions between sublevels of the $S_{1/2}$ and $P_{1/2}$ states of the ^{6}Li atom.

5. CONCLUSION

Let us consider in more detail the basic requirements that must be met by the atoms and the fields in order to obtain polarization of the ion nuclei in resonance ionization of atoms. In addition to the requirements for selective excitations of hfs states generalized in condition (1) and the requirements on the polarizations of the fields discussed in Section 4, we must obviously also formulate conditions on the strengths of the fields.

Concerning the choice of the atoms and the resonant states, the following necessary requirement follows from what has been said above: the hfs must be resolved in one or both of the resonant states (the ground state and/or the excited state). For multielectron atoms one can only point to individual typical cases. Thus, for alkali atoms the hfs is resolved in both the ground state and the first excited state[12]; for the hydrogen atom, as was mentioned above, the hfs is resolved only in the ground state[9]; and one can point to atoms for which the reverse situation obtains—this is the case when the nuclear spin is zero, as in beryllium, for example.

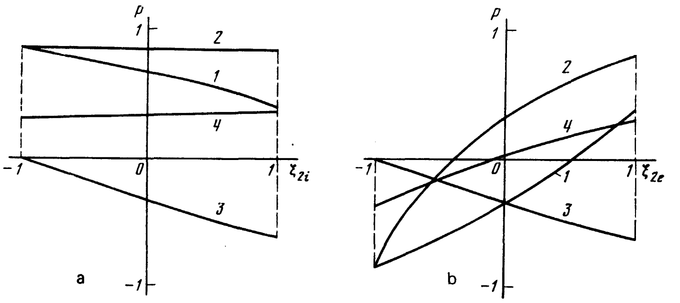

FIG. 3. Polarization P of the nuclei of ^{6}Li$^+$ ions vs the circular polarization ξ_{2i} of the ionizing radiation with $\xi_{2e}=1$ (plot a) and the circular polarization ξ_{2e} of the exciting radiation with $\xi_{2i}=1$ (plot b). The total angular momentum quantum numbers F and F' of the ground and excited states have the following values for the several curves: 1—$F=F'=\frac{1}{2}$; 2—$F=\frac{1}{2}$, $F'=\frac{3}{2}$; 3—$F=\frac{3}{2}$, $F'=\frac{1}{2}$; 4—$F=F'=\frac{3}{2}$.

The high spectral purity of the laser light required by condition (1) ($\Delta>\Delta\omega$) is within the capabilities of modern lasers operating in a single mode with fixed transverse and longitudinal indices. Correspondingly, there is a lower bound on the duration of the laser pulse: $\tau\gtrsim 1/\Delta\omega > 1/\Delta \sim 10^{-11}-10^{-8}$ sec. This range is typical for Q-switched pulsed lasers. In classical language, the condition $\tau>1/\Delta$ means that the duration of the light-field pulse must exceed the precession period of the vectors j and I.

The problem of the optimal strengths of the exciting and ionizing fields was discussed in detail in Ref. 13 in connection with the polarization of the electrons produced in resoance ionization of atoms. All the conditions formulated in Ref. 13 for bringing about resonance ionization via an hfs state are valid for the case of nuclear polarization. Summarizing the results of Ref. 13, we can say that the optimum strength of the exciting and ionizing fields is determined by the condition $\Gamma(\mathscr{E})\sim\gamma$, where γ is the natural width of the hfs level and $\Gamma(\mathscr{E})$ is the width of that level in the presence of the external field. The polarization falls off when $\Gamma(\mathscr{E})>\gamma$. When the fieldstrength is so high that $\Gamma(\mathscr{E})\gtrsim\Delta(\mathscr{E})$ there will obviously be no polarization at all since then condition (1) for selective excitation will no longer be satisfied. Here $\Delta(\mathscr{E})$ is the characteristic hfs splitting energy with allowance for the change of the level energies in the presence of the external field.

Let us make some numerical estimates of the optimum fieldstrengths for a typical case of two-photon ionization in the presence of a one-photon resonance. The exciting field will lead to mixing of the resonant states and to the appearance of the so-called field width $\Gamma(\mathscr{E})\sim d_{01}\mathscr{E}$. For the typical values $\gamma\sim10^{-3}$ cm^{-1} and $d_{01}\sim 1$ D the condition $\Gamma(\mathscr{E})\sim\gamma$ first obtains when $\mathscr{E}\sim10^2$ V/cm. It should be noted that this estimate is quite well confirmed quantitatively by the experimental data on the yield and polarization of electrons from two-photon resonance ionization of the cesium atom via hfs states[6,7]: the electron yield begins to grow less rapidly than $\mathscr{E}^2$ and the polarization falls off when the strength of the exciting field exceeds ~50 V/cm. The ionizing field will increase the width of the resonance state because it increases the ionization probability, i.e., the probability that the electron will make a transition to the continuous spectrum. For the ionization broadening that arises in this manner we have $\Gamma(\mathscr{E}')\sim w_{1E}\sim(\mathscr{E}')^2$; for the typical value $\sigma\sim10^{-17}$ cm^2 for the photoionization cross section it becomes of the order of the natural width of the level at the field strength $\mathscr{E}\sim10^5$ V/cm. In particular, it is evident from the above data that to optimize the yield of polarized nuclei one must have two distinct fields: an exciting field and an ionizing field.

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"Destruction" of hydrogen atom by collisions with multiply charged ions

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A calculation is reported of the total cross section of "destruction" of a hydrogen atom by collision with a multiply charged ion ($Z \gg 1$). The partial cross sections representing charge exchange between a hydrogen atom and a multiply charged ion are calculated using perturbation theory for collision velocities low compared with $e^2/\hbar$ (e is the electron charge and $\hbar$ is the Planck constant). The total charge-exchange cross section is compared with the results of a model in which the number of final states of a multiply charged ion is regarded as infinite. The criterion of validity of this model is determined. The coordinate and time dependences of the argument of the exponential electron wave function are determined for collision velocities low compared with $Z^{1/2}e^2/\hbar$ (Z is the ion charge) using the quasiclassical Keldysh method. The ionization cross section of a hydrogen atom colliding with a multiply charged ion is calculated for collision velocities high compared with $Z^{1/2}e^2/\hbar$. Matching of these cross sections makes it possible to determine the total destruction cross section of a hydrogen atom colliding with a multiply charged ion, which is valid in a wide range of collision velocities.

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The parameters of a thermonuclear (fusion) plasma and the rates of heating and decay of such a plasma are affected considerably by impurities.[1,2] In a fusion plasma the impurities are present in the form of multiply charged ions so that their influence on the plasma properties if manifested in various processes involving such ions. In particular, when a beam of fast hydrogen atoms is injected into a plasma, the processes of "destruction" of hydrogen atoms by collisions with multiply charged ions are important. The present paper is concerned with a calculation of the cross sections of such processes.

The cross sections of inelastic processes in collisions of multiply charged ions with atoms or singly charged ions have been determined in many recent experiments[3-8] and found in many calculations.[9-19] A characteristic of these processes is associated with a large number of possible reaction channels because there are many electron states in the field of a multiply charged ion to which the electron can be transferred. The existing theoretical approaches to the calculation of cross sections of such processes are based on the two-level approximation or on perturbation theory and are designed to determine the partial cross sections of the investigated processes. They require detailed information on the spectrum of multiply charged ions, which can be obtained only after laborious calculations, and the results apply to specific partners and parameters of collisions. Such calculations cannot always be applied to other collision parameters and other partners.

Our aim is to find the total cross section for the process of an inelastic collision between a hydrogen atom and a multiply charged ion. The loss of information on the details of the process makes it possible to combine all the channels that "destroy" the hydrogen atom.

We shall use asymptotic and quasiclassical methods. We shall divide the range of collision velocities arbitrarily into three regions. In the first region the relative collision velocity is small compared with the character-

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