

# The Fano effect induced by a strong external electromagnetic field

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A method is proposed for obtaining polarized electrons by ionizing atoms in two external fields having different frequencies  $\omega_0$  and  $\omega_1$  that satisfy the condition  $E_0 + \omega_0 \approx E_1 + \omega_1$  for resonance via the continuum of two discrete levels with energies  $E_0$  and  $E_1$ . The analogy between this phenomenon and the familiar Fano effect is traced. The physical causes of the polarization—the spin-orbit coupling and level splitting due to the dynamic Stark effect—are discussed. It is shown that under optimal conditions the electron polarization may be very high and that the proposed method is quite a universal one. The method can be used with atoms of virtually any type, with any level  $E_1$  in resonance with the ground state  $E_0$ , and over a broad frequency range (since the level coupling the ground state  $E_0$  with the level  $E_1$  may lie anywhere in the continuum).

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## 1. INTRODUCTION. FORMULATION OF THE PROBLEM

The spin of the electron is of vital importance in many physical problems involving collisions between electrons and atoms, or nuclear physics. Hence the construction of sources of polarized electrons, i.e., of electrons having their spins preferentially oriented in some specific direction, is of great interest. Quite a few different physical principles on which methods for constructing polarized-electron sources could be based are now known. Of these methods, those involving the interaction of atoms with electromagnetic radiation<sup>1</sup> form a special group, of which the method based on the Fano effect<sup>2</sup> would seem to be the first to have been proposed. As is well known, the Fano effect takes place in atoms whose photoionization cross section has a deep minimum (the so-called Cooper minimum<sup>3</sup>) as a function of the kinetic energy of the photoelectron. The spin-orbit interaction plays a relatively more important part in the vicinity of the minimum, and as a result the electrons produced by photoionization of the atom by circularly polarized radiation may be very highly (up to 100%) polarized (Fig. 1).

By now the Fano effect has been observed and fairly thoroughly investigated experimentally<sup>4,5</sup> (also see the references cited in the review article by Delone and one of us<sup>1</sup>). Despite all the advantages of the Fano effect, this method also has some advantages. First, since considerable polarization of the electrons is achieved only in the vicinity of the minimum in the photoionization cross section, the absolute photoelectron yield will

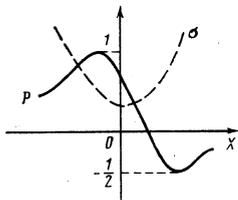


FIG. 1. Ionization cross section  $\sigma$  and photoelectron polarization  $P$  near the Cooper minimum;  $x \sim E_0 - E$ , where  $E_0$  is the kinetic energy at which the cross section is minimum.

be comparatively low (in other words, more intense electromagnetic radiation will be required for a given photoelectron yield than would be necessary far from the Cooper minimum). Another disadvantage of the method is that the Fano effect is far from universal: not all atoms exhibit the Fano effect. Finally, we note that the condition for photoionization in the vicinity of the Cooper minimum strictly fixed the frequency of the ionizing radiation, and the required frequency may turn out to be quite unsuitable.

As will be shown below, there is a photoionization scheme that partly or completely avoids these disadvantages. The proposed method is based on the idea that what is important for the presence of the Fano effect is essentially only that the photoionization cross section have a sufficiently deep minimum, but that the mechanism giving rise to that minimum does not matter much. A scheme for the photoionization of atoms by the action of two fields having different frequencies and intensities, which satisfies this condition, is now known and has been fairly well investigated theoretically (Fig. 2).<sup>6-8</sup> Let  $E_0$  and  $E_1$  be the energies of the ground state and an excited S state of an atom, and let the frequencies  $\omega_0$  and  $\omega_1$  of electromagnetic fields with the field strengths  $\text{Re}(\mathbf{F}_0 \exp(-i\omega_0 t))$  and  $\text{Re}(\mathbf{F}_1 \exp(-i\omega_1 t))$  satisfy the condition  $E_1 + \omega_1 \approx E_0 + \omega_0$  for resonance via the continuum. Finally, let the field  $\mathbf{F}_1$  be strong and the field  $\mathbf{F}_0$ , weak. Then if we neglect the spin-orbit coupling we find that the photoionization probability, regarded as a function of the frequency  $\omega_0$  of the weak field, will be characterized

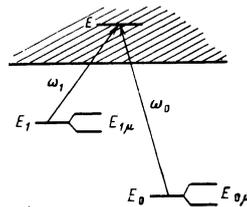


FIG. 2. Scheme for the ionization of an atom in two fields of frequencies  $\omega_0$  and  $\omega_1$  satisfying the condition  $E_1 + \omega_1 \approx E_0 + \omega_0$  for resonance via the continuum. The figure also shows the splitting of the S levels  $E_0$  and  $E_1$  by the dynamic Stark effect.

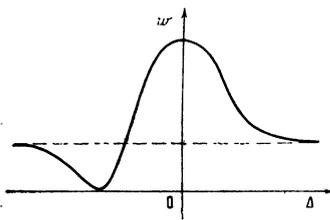


FIG. 3. Ionization probability  $w$  near the resonance as a function of the detuning  $\Delta \approx E_1 + \omega_1 - E_0 - \omega_0$ .

by an asymmetric curve like the one sketched in Fig. 3 and will vanish for some value of  $\omega_0$  near the resonance. The polarization of the electrons produced in such an ionization scheme has not previously been investigated. On the basis of the considerations advanced above we may assume that when the spin-orbit coupling is taken into account the minimum ionization probability in this scheme will be found to be finite and that, in analogy with the Fano effect, considerable polarization of the photoelectrons produced in the vicinity of the minimum will be possible. These assumptions are fully confirmed by the calculations presented below.

In the ionization scheme under consideration there is, apart from the spin-orbit coupling, another physical factor that could result in a finite minimum value for the ionization probability and in a comparatively high degree of polarization of the electrons. This factor is the possible splitting of the levels  $E_0$  and  $E_1$  into sublevels  $E_{0\mu}$  and  $E_{1\mu}$  corresponding to the two projections  $\mu = \pm 1/2$  of the electron spin onto the  $z$  axis as a result of the dynamic Stark effect in the strong external field  $F_1$ . We note that the polarization of electrons as a result of Stark splitting of atomic levels in a different ionization scheme, namely in the resonant multiphoton ionization of atoms, has been discussed theoretically in Refs. 9 and 10.

For definiteness we shall assume that the two waves  $F_0$  and  $F_1$  propagate in the same direction (along the  $z$  axis) and are both right-hand circularly polarized, and for simplicity we shall analyze the transitions in atoms having but a single valence electron.

## 2. QUASIENERGY ATOMIC WAVE FUNCTIONS IN THE STRONG FIELD

Let us first consider the effect of the strong field  $F_1$  on the spectrum and wave functions of the atom. It is obvious that the external field of frequency  $\omega_1$ , whose frequency is higher than the frequency corresponding to the ionization potential of the level  $E_1$ , will mix the states of this level with states of the continuum. It is known<sup>9</sup> that as a result of such mixing there arise quasienergy states with a continuous spectrum of quasienergies  $E$ . In the present case the problem of finding the quasienergy wave functions is closely analogous to the corresponding problem that arises in treating the direct ionization of an atom by a strong external field,<sup>11</sup> as well as to the steady state problem of discrete levels against a continuous background.<sup>12</sup> A special feature of the present case is that the electron spin and the spin-orbit coupling are to be taken into account. Each of the levels  $E_0$  and  $E_1$  must therefore be regarded as a pair of sublevels corresponding to the two projections  $\mu = \pm 1/2$

of the spin onto the  $z$  axis.

The strong external field  $F_1$  not only mixes the states of the discrete spectrum with continuum states, but also shifts the levels  $E_0$  and  $E_1$  and splits them as a result of the quadratic dynamic Stark effect:

$$E_0 \rightarrow E_{0\mu} = E_0 - 1/4 \alpha_{0\mu} F_1^2, \quad E_1 \rightarrow E_{1\mu} = E_1 - 1/4 \alpha_{1\mu} F_1^2, \quad (1)$$

where  $\alpha_{0\mu}$  and  $\alpha_{1\mu}$  are the dynamic polarizabilities of the levels. In what follows we shall assume that these shifts and splittings have been already taken into account in the energies  $E_0$  and  $E_1$ . The part of the shift of the levels  $E_{1\mu}$  due to transitions via the continuum, which arises explicitly in the subsequent calculations, must therefore be omitted.

The operator  $V = -(1/2)d \cdot F_1$  for the interaction of the atom with the field  $F_1$  ( $d$  is the dipole moment) couples the states  $|1S_{1/2, \mu}\rangle$  with the  $P$  states  $|EP_{j, \mu+1}\rangle$  of the continuous spectrum ( $E$  is the energy,  $j = \frac{1}{2}, \frac{3}{2}$  is the total angular momentum of the electron, and  $\mu + 1 \leq j$  is the projection of the electron angular momentum onto the  $z$  axis). It is convenient to introduce the following linear combinations of continuum wave functions:

$$\begin{aligned} \psi_{E-\nu_2} &= \frac{1}{v_-(E)} \sum_j \left| EP_{j, \frac{1}{2}} \right\rangle v_j(E), \\ \psi_{E0} &= \frac{1}{v_-(E)} \left\{ \left| EP_{\nu_2, \frac{1}{2}} \right\rangle v_{\nu_2}^*(E) - \left| EP_{\nu_2, \frac{1}{2}} \right\rangle v_{\nu_2}(E) \right\}, \\ \psi_{E\nu_2} &= \left| EP_{\nu_2, \frac{3}{2}} \right\rangle; \end{aligned} \quad (2)$$

in which

$$v_j(E) = \left\langle EP_{j, \frac{1}{2}} \left| V \right| 1S_{\nu_2, -\frac{1}{2}} \right\rangle, \quad v_-(E) = \left( \sum_j |v_j(E)|^2 \right)^{1/2}. \quad (3)$$

The functions  $\psi_{E-1/2}$  and  $\psi_{E1/2}$  are coupled to the states  $|1S_{1/2, -1/2}\rangle$  and  $|1S_{1/2, 1/2}\rangle$ , respectively. The states represented by the wave functions  $\psi_0$  do not mix with the states  $|1S_{1/2, 2\mu}\rangle$  to its own continuum, whose wave functions are  $\psi_{E\mu}$ . Now using the above mentioned analogy with the work of Kazakov et al.<sup>11</sup> we can immediately write out explicitly the complete set of quasienergy functions for the atom in the strong field  $F_1$ , with quasienergy  $E$  (we use units in which  $\hbar = 1$ ):

$$\begin{aligned} \Psi_{E\pm\nu_2} &= \frac{e^{-iEt}}{v_{\pm}(E) (z_{\pm}^2(E) + \pi^2)^{1/2}} \left\{ e^{i\omega_1 t} \left| 1S_{\nu_2, \pm \frac{1}{2}} \right\rangle \right. \\ &+ \int dE' v_{\pm}(E') \Psi_{E', \pm\nu_2} \left( P \frac{1}{E - E'} + z_{\pm}(E) \delta(E' - E) \right) \left. \right\}, \\ \Psi_{E0} &= e^{-iEt} \psi_{E0}, \end{aligned} \quad (4)$$

where  $v_{\pm}(E)$  is defined by Eqs. (3), and

$$v_{\pm}(E) = \langle \psi_{E\nu_2} | V | 1S_{\nu_2, \pm 1/2} \rangle = \langle EP_{\nu_2, \pm 1/2} | V | 1S_{\nu_2, \pm 1/2} \rangle. \quad (5)$$

The principle-value integrals in formulas (4) and everywhere in what follows are to be understood as integrals over the continuous spectrum plus summations over all the nonresonant states of the discrete spectrum. The functions  $z_{\pm}(E)$  have the form

$$z_{\pm}(E) = (E - E_{1, \pm\nu_2} - \omega_1) / (v_{\pm}(E))^2. \quad (6)$$

In accordance with the assumption, noted above, that the Stark shift is fully taken into account in the energies

$E_{1\mu}$ , the principle-value integrals that determine the part of this shift due to transitions via the continuum (cf. formula (6) in Ref. 11) are omitted from formula (6).

### 3. IONIZATION PROBABILITY AND ELECTRON POLARIZATION. GENERAL FORMULAS

Now let us use perturbation theory to take account of the transitions from the ground state to the quasienergy continuum states (4) induced by interaction with the field  $F_0$ . The matrix elements of the interaction operator  $U = -(1/2)d \cdot F_0$  between the states  $|0S_{1/2, \mu}\rangle$  corresponding to the sublevels of the state  $E_{0\mu}$  and the continuum functions  $\psi_{E \mp 1/2}$  have the form

$$u_-(E) = \langle \psi_{E-1/2} | U | 0S_{1/2} - \frac{1}{2} \rangle = \frac{1}{v_-(E)} \sum_j v_j^*(E) u_j(E), \quad (7)$$

$$u_+(E) = \langle \psi_{E1/2} | U | 0S_{1/2, 1/2} \rangle = \langle EP_{3/2}^+ | U | 0S_{1/2, 1/2} \rangle, \quad (8)$$

$$u_j(E) = \langle EP_{j1/2}^+ | U | 0S_{1/2, 1/2} \rangle.$$

In addition, the matrix element for the transition to the state  $\psi_{E0}$  is also different from zero:

$$u_0(E) = \langle \psi_{E0} | U | 0S_{1/2} - \frac{1}{2} \rangle = \frac{u_{1/2}(E)v_{1/2}(E) - u_{3/2}(E)v_{3/2}(E)}{v_-(E)}. \quad (9)$$

We shall assume that the interaction with the field  $F_0$  is turned on suddenly at time  $t = 0$ . The solution  $\Psi(t)$  to the initial-value problem depends essentially on which of the states  $|0S_{1/2, \mu}\rangle$  the atom was in at  $t = 0$ . In the two cases

$$\Psi(0) = |0S_{1/2, 1/2}\rangle \text{ and } \Psi(0) = |0S_{1/2, 3/2}\rangle$$

the expansion of the wave functions  $\Psi_{\mp 1/2}(t)$  in the complete set of quasienergy atomic functions (4) has the form

$$\Psi_{-1/2}(t) = \exp(-iE_{0-1/2}t) \left| 0S_{1/2} - \frac{1}{2} \right\rangle + \int dE C_{E-1/2}(t) \Psi_{E-1/2} + \int dE C_{E0}(t) \Psi_{E0}, \quad (10)$$

$$\Psi_{1/2}(t) = \exp(-iE_{01/2}t) \left| 0S_{1/2, 1/2} \right\rangle + \int dE C_{E1/2}(t) \Psi_{E1/2}. \quad (11)$$

In first order perturbation theory, the expansion coefficients  $C_{E\mu}(t)$  and  $C_{E0}(t)$  for large  $t$  can be written in the form

$$C_{E\mp 1/2} = -\frac{2\pi i \delta(E - \omega_0 - E_{0\mp 1/2})}{v_{\mp}^2(E) (z_{\mp}^2(E) + \pi^2)} \int dE' v_{\mp}^*(E') \times u_{\mp}(E') \left( P \frac{1}{E - E'} + z_{\mp}(E) \delta(E - E') \right), \quad (12)$$

$$C_{E0} = -2\pi i u_0(E) \delta(E - E_{01/2} - \omega_0). \quad (13)$$

In calculating the ionization probability and the polarization of the photoelectrons we shall assume that the atoms were initially unpolarized, i.e., that they were in the states  $|0S_{1/2, 1/2}\rangle$  and  $|0S_{1/2, 3/2}\rangle$  with equal probability. All physical quantities must accordingly be averaged over these initial states. The probability for ionizing the atom is determined by projecting the functions  $\Psi_{\mu}(t)$  [Eqs. (10) and (11)] onto continuum states of the free atom and then integrating the squares of the transition probability amplitudes over the energy:

$$w(t) = \frac{1}{2} \int dE \{ |\langle \psi_{E-1/2} | \Psi_{-1/2}(t) \rangle|^2 + |\langle \psi_{E0} | \Psi_{-1/2}(t) \rangle|^2 + |\langle \psi_{E1/2} | \Psi_{1/2}(t) \rangle|^2 \}. \quad (14)$$

Without pausing to discuss the calculation of the integrals that determine  $w(t)$ , we give the final expression for the ionization probability in the large- $t$  limit where the probability is proportional to  $t$ :

$$w = \pi t \{ |u_0(E_{0-1/2} + \omega_0)|^2 + \sum_{\pm} v_{\pm}^{-2}(E_{0\pm 1/2} + \omega_0) [z_{\pm}^2(E_{0\pm 1/2} + \omega_0) + \pi^2]^{-1} \times \left| \int dE v_{\pm}^*(E) u_{\pm}(E) \left( P \frac{1}{E_{0\pm 1/2} + \omega_0 - E} + z_{\pm}(E) \delta(E_{0\pm 1/2} + \omega_0 - E) \right) \right|^2 \}. \quad (15)$$

The degree of polarization  $P$  is defined as twice the ratio of the average projection  $S_z$  of the photoelectron onto the  $z$  axis to the ionization probability  $w$ ; calculations analogous to those presented above give

$$P = \frac{\langle 2S_z \rangle}{w} = \frac{2\pi t}{w} \left\{ \langle \psi_{E0} | S_z | \psi_{E0} \rangle |u_0(E_{0-1/2} + \omega_0)|^2 + \sum_{\pm} \frac{\langle \psi_{E\pm 1/2} | S_z | \psi_{E\pm 1/2} \rangle}{v_{\pm}^2(E_{0\pm 1/2} + \omega_0) (z_{\pm}^2(E_{0\pm 1/2} + \omega_0) + \pi^2)} \left| \int dE v_{\pm}^*(E) u_{\pm}(E) \times \left( P \frac{1}{E_{0\pm 1/2} + \omega_0 - E} + z_{\pm}(E) \delta(E_{0\pm 1/2} + \omega_0 - E) \right) \right|^2 \right\}. \quad (16)$$

The wave functions between which the matrix elements in formulas (15) and (16) are calculated are products of spherical functions by radial wave functions for the valence electron. The spin-orbit coupling obviously affects only the radial functions and the radial matrix elements; this makes it possible to push the ionization-probability and polarization calculations somewhat further while retaining the general form.

Let the radial matrix elements for the transitions  $|1S_{1/2, \mu}\rangle \rightarrow |EP_{j, \mu} + 1\rangle$  and  $|0S_{1/2, \mu}\rangle \rightarrow |EP_{j, \mu} + 1\rangle$  be denoted by  $R_{1,3}^{(1)}$  and  $R_{1,3}^{(0)}$ , respectively, where the lower indices 1 and 3 correspond to  $j = 1/2$  and  $j = 3/2$ , respectively. We also write

$$Q = [(R_3^{(1)})^2 + 2(R_1^{(1)})^2]^{1/2}.$$

The difference between  $R_1^{(\alpha)}$  and  $R_3^{(\alpha)}$  is due to the effect of the spin-orbit coupling on the matrix elements for transitions from the levels  $E_0$  and  $E_1$  to the continuum. (If the spin-orbit coupling were neglected we would have  $R_1^{(\alpha)} = R_3^{(\alpha)} = R^{(\alpha)}$  and  $Q = 3^{1/2} R^{(1)}$ .)

On calculating the angular parts of the matrix elements and substituting them into formulas (15) and (16), these formulas take the form

$$w = \pi t e^2 F_0^2 \{ C + A(\Delta_{-1/2} + a)^2 + B(\Delta_{1/2} + b)^2 \}, \quad (17)$$

$$P = \frac{2\pi t e^2 F_0^2}{w} \left\{ \frac{2(R_1^{(1)})^2 + 8R_1^{(1)}R_3^{(1)} - (R_3^{(1)})^2}{6Q^2} C + \frac{(R_3^{(1)})^2 - 8R_1^{(1)}R_3^{(1)} - 2(R_1^{(1)})^2}{6Q^2} A(\Delta_{-1/2} + a)^2 + \frac{1}{2} B(\Delta_{1/2} + b)^2 \right\}, \quad (18)$$

where  $\Delta_{\mu} = E_{1\mu} - E_{0\mu} + \omega_1 - \omega_0$  is the deviation from resonance,

$$C = (R_3^{(1)}R_1^{(0)} - R_1^{(1)}R_3^{(0)})^2/Q^2, \quad (19)$$

$$A = \frac{1}{18Q^2} \frac{(2R_1^{(1)}R_1^{(0)} + R_3^{(1)}R_3^{(0)})^2}{\Delta_{-1/2}^2 + \pi^2 e^4 F_1^4 Q^2/324}, \quad (20)$$

$$a = \frac{e^2 F_1^2 Q^2}{18[2R_1^{(1)}R_1^{(0)} + R_3^{(1)}R_3^{(0)}]} \int dE \frac{2R_1^{(1)}(E)R_1^{(0)}(E) + R_3^{(1)}(E)R_3^{(0)}(E)}{E_{0-1/2} + \omega_0 - E} \quad (21)$$

$$B = \frac{1}{6} \frac{(R_3^{(0)})^2}{\Delta_{1/2} + \pi^2 e^4 F_1^4 (R_3^{(1)})^2/36}, \quad (22)$$

and

$$b = \frac{e^2 F_1^2 R_3^{(1)}}{6 R_3^{(0)}} \int dE \frac{R_3^{(1)}(E) R_3^{(0)}(E)}{E_{01/2} + \omega_0 - E}. \quad (23)$$

The energy arguments of some of the radial matrix elements have been omitted; the missing arguments should be  $E_{0-1/2} + \omega_0$  in formulas (18)–(21), and  $E_{01/2} + \omega_0$  in formulas (22) and (23).

#### 4. ELECTRON POLARIZATION NEAR THE TRANSITION-PROBABILITY MINIMUM

Formulas (17) and (18) are fairly cumbersome and difficult to analyze in the general case. We shall therefore consider the most characteristic special cases.

1. If we neglect the spin-orbit interaction and the Stark splitting of the levels  $E_0$  and  $E_1$ , we have  $C = 0$ ,  $A = B$ ,  $a = b$ , and

$$\Delta_n \approx \Delta = \frac{1}{2} \sum_n \Delta_n \approx E_1 - E_0 + \omega_1 - \omega_0.$$

Then the polarization vanishes ( $P = 0$ ) and Eq. (17) for the probability  $w$  assumes the usual<sup>6-8</sup> form

$$w = \frac{\pi t e^2 F_0^2 (R_3^{(0)})^2 (\Delta + a)^2}{3 (\Delta^2 + \pi^2 e^2 F_1^4 (R^{(1)})^4 / 36)}. \quad (24)$$

Formula (24) corresponds to the curve shown in Fig.

3. When  $\Delta = -a$  we have  $w = 0$ ; the shift of this point from the point of exact resonance ( $\Delta = 0$ ) is  $-a \sim -E_{at} (F_1 / F_{at})^2$ .

Now let us include the spin-orbit coupling and the Stark splitting of the levels  $E_0$  and  $E_1$ . Moreover, we shall assume that the Stark splitting, characterized by the parameter  $\Delta_0 \equiv \Delta_{1/2} - \Delta_{-1/2}$ , is small as compared with the constants  $a$  and  $b$ . This condition is usually satisfied in view of the fact that the coefficients of  $F_1^2$  in formulas (21) and (23) for  $a$  and  $b$  are of the same order as the ordinary scalar polarizabilities of the levels  $E_0$  and  $E_1$ . At the same time, in the case of S states the difference  $\Delta_0$  between the detunings is determined by the so-called vector polarizabilities of the levels, and these, as a rule, are one or two orders of magnitude smaller than the scalar polarizabilities.<sup>10</sup>

Since the constants  $a$  and  $b$  also differ little from one another (they differ only because of the spin-orbit coupling) the second and third terms in formula (17), regarded as functions of the average detuning  $\Delta$ , vanish at the points  $\Delta_0/2 - a$  and  $-\Delta_0/2 - b$ ; these points are close together, their separation being considerable smaller than their average shift from the position of the ionization-probability maximum. If one is interested only in comparatively small scale changes of the detuning of the order of  $\Delta_0 + b - a$ , one can simplify formulas (17) and (18) considerably by retaining only the lower order terms in the spin-orbit coupling and in  $\Delta_0$ . In this approximation the ionization probability assumes the form

$$w = \pi t e^2 F_0^2 (C + 1/2 A_0 (\Delta_0 + b - a)^2 + 2 A_0 x^2), \quad (25)$$

where

$$x = \Delta + \frac{a+b}{2} = \frac{1}{2} \left( \sum_n \Delta_n + a + b \right), \quad (26)$$

$$A_0 = \frac{(R^{(0)})^2}{1/2 (a+b)^2 + 1/6 \pi^2 e^2 F_0^4 (R^{(1)})^2}.$$

Formula (25) describes the parabolic change in the ionization probability  $w$  within a small neighborhood of its minimum value.

In this approximation the polarization, determined by formula (18), becomes

$$P = \frac{C + 2 A_0 (\Delta_0 + b - a) x}{C + 1/2 A_0 (\Delta_0 + b - a)^2 + 2 A_0 x^2}. \quad (27)$$

The polarization reaches its extreme values at the points

$$x_1 = \frac{\Delta_0 + b - a}{2}, \quad x_2 = -\frac{C + 1/2 A_0 (\Delta_0 + b - a)^2}{A_0 (\Delta_0 + b - a)}. \quad (28)$$

the corresponding extreme values being

$$P(x_1) = 1, \quad P(x_2) = -\left[ 1 + \frac{2C}{A_0 (\Delta_0 + b - a)^2} \right]^{-1}. \quad (29)$$

When  $\Delta_0 + b - a > 0$ , the function  $P(x)$  defined by Eq. (27), which is shown graphically in Fig. 4, is qualitatively similar to the frequency dependence of the electron polarization in the ordinary Fano effect (Fig. 1), the only difference being that in the present case (but not in the Fano effect)  $P_{\min} \neq 1/2$ . The maximum polarization achievable with the present scheme is 100%, just as in the case of the Fano effect.

3. Formula (27) is general in the sense that it takes into account both the electron polarization mechanisms: the spin-orbit coupling and the Stark splitting of the levels. The relation between these mechanisms is determined by the relation between  $\Delta_0$  and the parameters  $b - a$  and  $(C/A_0)^{1/2}$ . In view of the fact that the relative magnitude of the vector polarizability of the levels may vary between rather wide limits, the relation between the effects associated with the spin-orbit coupling and with the Stark splitting may differ in different atoms, for different levels, and at different frequencies. We shall therefore consider the case in which the Stark splitting of the levels is the principal polarization mechanism, assuming that

$$b - a \gg |\Delta_0| \gg |b - a| \sim (C/A_0)^{1/2}.$$

In this case Eq. (27) simplifies, assuming the form

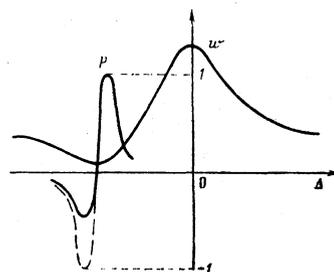


FIG. 4. Ionization probability and polarization with allowance for the spin-orbit interaction and the Stark splitting of the levels. The dashed curve is for the case in which the Stark splitting predominates (formula (30)).

$$P = \frac{\Delta_0 x}{x^2 + 1/\Delta_0^2}. \quad (30)$$

A special feature of this case is the symmetry of the  $P(x)$  curve. The extreme values of  $P(x)$  are reached at the points  $x_{1,2} = \pm \Delta_0/2$ , the extreme values being  $P_{\max, \min} = P(\pm \Delta_0/2) = \pm 1$ .

A rigorous quantitative determination of the conditions under which one or the other of the polarization mechanisms is predominant for any specific atom would require accurate calculation of the parameters  $C$ ,  $a$ ,  $b$ , and  $A_0$  defined by formulas (19), (21), (23), and (26). In principle, of course, such calculations could be made, but at present no results of such calculations are to be found in the literature. We are therefore forced to limit ourselves here to the general formulas (27) and (30) and to the illustrative graph of the function  $P(x)$  given in Fig. 4. It is important to note, however, that  $|\Delta_0| \ll a$  and  $|a - b| \ll a$ , as is evident from the general considerations presented above, and that one or the other of the electron-polarization mechanisms can therefore be realized with virtually any atom.

## 5. CONCLUSION

Thus, the scheme for ionizing atoms considered above makes it possible to obtain polarized electron beams with polarizations up to 100%. We would call attention to the fact that the method is highly universal. Indeed, the method can be used with virtually any atoms. All that is necessary is that the photoionization in the field of frequency  $\omega_0$  take place in the presence of intense radiation  $F_1$  at the "resonance" frequency  $\omega_1 \approx E_0 + \omega_0 - E_1$ . There are virtually no limitations on the choice of the level  $E_1$  except such as may be imposed for convenience. Since the resonance takes place via the continuum and the continuum level involved in the resonance can lie anywhere in the continuous spectrum, the frequency  $\omega_0$  and the frequency  $\omega_1$  associated with it can also vary between very wide limits. Thus, the intense external field  $F_1$  makes it possible to realize a phenomenon analogous to the Fano effect under very diverse conditions.

The criterion determining how strong the external field  $F_1$  must be is that the natural width  $\Gamma_r$  of the level  $E_1$  (which has not yet been taken into account at all) be small as compared with the characteristic scale of the variations in the degree of polarization  $P(\Delta)$  (Fig. 4). Let us estimate the corresponding limitations on the strength of the field  $F_1$ , assuming for simplicity that  $|\Delta_0| \sim |b - a| \sim (C/A_0)^{1/2}$ . Since we obviously have

$(C/A_0)^{1/2} \sim \alpha E_{at} (F_1/F_{at})^2$ , where  $\alpha$  is the fine structure constant (see Eqs. (19) and (26)), the condition  $|\Delta_0| > \Gamma_r$  gives  $F_1 > 5 \times 10^6$  V/cm provided  $\Gamma_r \sim 10^{-3}$  cm $^{-1}$ . It should be noted that the widths  $\Gamma_r$  of high-lying levels may actually be much smaller than  $10^{-3}$  cm $^{-1}$ , and that the limitation on  $F_1$  may therefore be less stringent. Moreover, the requirement even as estimated above is acceptable at the current level of laser development.

We also note that at first glance it might seem that the quantity  $\Delta_0$  must be compared not only with the natural line width  $\Gamma_r$ , but also with the ionization widths of the levels  $E_{1\mu}$ . Actually, this is not necessary since all the effects associated with ionization are automatically taken accurately into account in the formalism. In itself, the structure of the function  $P(\Delta)$  graphed in Fig. 4 is a result of the interference of different ionization channels (ionization directly from the level  $E_0$  and via the level  $E_1$ ).

In summing up we emphasize that the proposed method is quite universal and simple. We feel that the experimental validation of the method could be of general physical interest, and could also be important from the point of view of applications, i.e., for producing efficient and convenient sources of polarized electrons.

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