Effect of resonant electromagnetic field on the autoionizing states of atoms

A. I. Andryushin, A. E. Kazakov, and M. V. Fedorov

P. N. Lebedev Physics Institute, USSR Academy of Sciences
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The effect of a strong electromagnetic field that is at resonance with the transition between certain autoionizing and discrete levels of the atom on the photoabsorption spectrum of a second weak (probing) pulse is investigated. The conditions are obtained under which one of the components of the dispersion curve can become narrower in a strong field. The minimum width of the narrow peak is obtained and the physical factors that determine the width are discussed.

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1. INTRODUCTION

Much attention is being paid of late to theoretical and experimental investigations of autoionizing states (AS) of atoms [1, 2]. These states can be realized in principle in any multielectron atom. Autoionizing states have a finite width governed both by the spontaneous emission of the photon (radiation width) and by the autoionization, i.e., the spontaneous decay into an ion and a free electron on account of the electrostatic interaction of the atomic electrons. Although the width corresponding to the autoionization channel varies in a wide range (10^{-1}–10^{4} cm^{-1}), it is usually much larger than the radiation width (~10^{-2} cm^{-1}) and it is therefore the deciding factor in the total width of the AS.

An investigation of the AS of multielectron atoms is of importance both for the study of the atomic physics proper, and for applications. For example, narrow AS (noble gases, alkaline-earth elements, lanthanides) are of interest in systems for selective multistep laser photoionization of atoms for isotope separation.[3] The excitation of such states with their subsequent autoionization can ensure a larger yield of the sought isotope than direct photoionization of the atoms.

In the presence of AS, the photoabsorption spectra have a resonant shape with a characteristic asymmetrical Fano profile: one wing of the dispersion coefficient has a deep minimum, this being the consequence of interference between the transitions of the atom first into the AS at large detuning from resonance and then into the continuum, on the one hand, and when the atom is directly ionized, on the other.[1,4] We note here that states similar to autoionizing can be induced in any atom by strong monochromatic electromagnetic field that leads to the onset of discrete quasi-energy states against the background of a continuum, if the field frequency is higher than the ionization threshold of the corresponding real levels.[5–10]

Laser (including multiphoton) spectroscopy methods are presently extensively used to investigate the characteristics of AS.[11,12] An investigation of the physics of the interaction of a strong electromagnetic field with AS can be of importance, in particular, for the interpretation of experiments on two-electron ionization of atoms of a second group.[13]

Analysis of the influence of laser radiation on AS is in fact only beginning at present.[14,15] In Ref. 14 it is indicated that a strong narrowing of the AS (to the radiative widths) is possible under the influence of electromagnetic radiation at resonance with the transition frequency between the given AS and a certain discrete level of the atom. It seems to us, however, that this result is due essentially to the model character of the problems considered in Ref. 14. Some of the assumptions made in these papers may actually not hold true. For example, it is always necessary to take into account transitions from AS to higher continuum states under the influence of a strong field. The problem is formulated more realistically in the next section. The solution of the problem will show that under certain conditions it is indeed possible to narrow down the dispersion curves. This, however, calls for a strictly defined selection of the field parameters. The minimum widths of the resonances are determined not by the radiative widths, but primarily by other factors—the strong-field-induced ionization of the autoionizing level, and the incomplete interference between the autoionization transition and the transition into the continuum from the AS via a discrete level under the influence of a strong field. These factors are disregarded also in other recent papers[15] devoted to the analysis of the influence of a strong field on the AS of an atom.

2. FORMULATION OF PROBLEM AND BASIC EQUATIONS

We consider resonant interaction of an AS with a certain discrete state of an atom in a strong field of frequency ω, $\hbar\omega \approx E_a - E_1$, where E_a and E_1 are the energies of the AS and of the discrete level. The strong field causes mixing of these states, i.e., formation of two quasi-energy states having certain widths both as a result of autoionization of the state E_a and as a result of the subsequent ionization of both states under the influence of the field. To investigate the obtained states we can use a weak (probing) field of frequency Ω, such that $\hbar\Omega \approx E_a - E_0$, where E_0 is the energy of the ground state of the atom. The weak field will be taken into account in the lowest order of perturbation theory, and the strong field will be taken into account exactly (in the resonance approximation). The transition scheme is shown in Fig. 1(a).

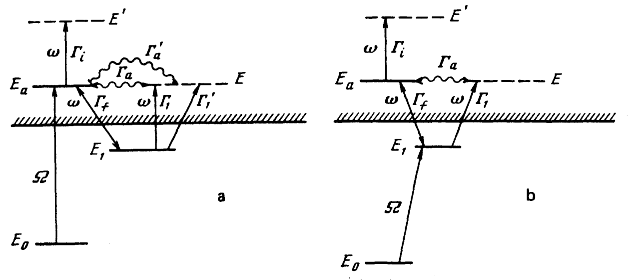

FIG. 1. Schemes of considered transitions.

We represent the wave function of the atom in the form of an expansion in unperturbed atomic functions of the ground state (ψ_0), the excited (ψ_1) and the autoionizing state (ψ_a), as well as of the states of the continuum ψ_E and ψ'_E with energies close respectively to E_a and $E_a+\hbar\omega$:

$$\Psi(t)=a_0(t)e^{i(\omega+\Omega)t}\psi_0+a_1(t)e^{2i\omega t}\psi_1+a_a(t)e^{i\omega t}\psi_a + e^{i\omega t}\int dE a_E(t)\psi_E+\int dE' a_{E'}(t)\psi_{E'}. \tag{1}$$

(Here and elsewhere we put $\hbar=1$.)

For simplicity we consider for the time being the case when all the states $\psi_0, \psi_1, \psi_a, \psi_E, \psi'_E$ are not degenerate. Generalizations to more complicated cases will be considered at the end of this section.

The operators of the interaction of the atom with the strong and with the probing fields are written in the form

$$V(t)=-{}^1/_2\mathbf{d}(\mathbf{E}e^{i\omega t}+\mathbf{E}^*e^{-i\omega t})\equiv Ve^{i\omega t}+V^+e^{-i\omega t},$$
$$F(t)=-{}^1/_2\mathbf{d}(\vec{\mathscr{E}}e^{i\Omega t}+\vec{\mathscr{E}}^*e^{-i\Omega t})\equiv Fe^{i\Omega t}+F^+e^{-i\Omega t},$$

where $\mathbf{E}$ and $\vec{\mathscr{E}}$ are the amplitudes of the corresponding electric fields, and $\mathbf{d}$ is the dipole-moment operator of the atom. We shall assume that both fields are turned on instantaneously at $t=0$, and that the unperturbed atom is in the ground state:

$$a_0(0)=1,\quad a_1(0)=a_a(0)=a_E(0)=a_{E'}(0)=0. \tag{2}$$

Substitution of the wave function (1) in the Schrödinger equation leads to a system of differential equations with constant coefficients for the amplitudes $a_0(t),\dots,$ $a_{E'}(t)$, which reduces, after taking the Laplace transform

$$a(p)=\int_0^\infty dt\, a(t)e^{-pt}$$

and allowing for the initial conditions (2), the following linear system of algebraic equations for the Laplace transforms $a(p)$ of the amplitudes $a(t)$:

$$\begin{gathered}[ip-(E_0+\Omega+\omega)]a_0(p)=i+F_{0a}a_a(p),\\ [ip-(E_1+2\omega)]a_1(p)=V_{1a}a_a(p)+\int dE\, V_{1E}a_E(p),\\ [ip-(E_a+\omega)]a_a(p)=F_{0a}^*a_0(p)+V_{1a}^*a_1(p)\\ +\int dE\, W_{aE}a_E(p)+\int dE'\, V_{aE'}a_{E'}(p),\\ [ip-(E+\omega)]a_E(p)=V_{1E}^*a_1(p)+W_{aE}^*a_a(p),\\ [ip-E']a_{E'}(p)=V_{aE'}^*a_a(p).\end{gathered} \tag{3}$$

The matrix elements of all the operators are calculated with the wave functions that enter in the expansion (1), and W_{aE} is the matrix element of the autoionization-decay operator. Transitions between the states of the continuous spectrum at $E\ll E_{at}=5\times10^9$ V/cm lead to small corrections (see, e.g., Ref. 16), and will therefore be disregarded. In Eqs. (3), we disregard also the direct transitions from the ground state into the continuum. This approximation is justified because our subsequent task is to describe the far wings of the resonance curves. It can be verified that in this case the interference of a direct and resonant transitions plays no role. If Γ is the characteristic width of some resonance maximum, then the interference between the direct and resonant transitions manifest itself upon detuning from the center of the maximum by an amount $\sim\sqrt{\Gamma}\gg\Gamma$ (in atomic units). These effects are not considered in the present paper.

In the lowest order in the weak field, $a_0^{(0)}(t)=1$ or

$$a_0^{(0)}(p)=\frac{i}{ip-(E_0+\omega+\Omega)}. \tag{4}$$

Eliminating with the aid of Eqs. (3) the amplitudes $a_E(p)$ and $a_{E'}(p)$, we obtain the following system of equations for the amplitudes $a_1^{(1)}(p)$ and $a_a^{(1)}(p)$ in the principal order in the weak field:

$$\begin{gathered}\left[ip-E_1-2\omega-\int dE\frac{|V_{1E}|^2}{ip-E-\omega}\right]a_1^{(1)}(p)-\left[V_{1a}+\int dE\frac{V_{1E}W_{aE}^*}{ip-E-\omega}\right]a_a^{(1)}(p)=0,\\ -\left[V_{1a}^*+\int dE\frac{W_{aE}V_{1E}^*}{ip-E-\omega}\right]a_1^{(1)}(p)+\left[ip-E_a-\omega\right.\\ \left.-\int dE\frac{|W_{aE}|^2}{ip-E-\omega}-\int dE'\frac{|V_{aE'}|^2}{ip-E'}\right]a_a^{(1)}(p)=F_{0a}^*a_0^{(0)}(p).\end{gathered} \tag{5}$$

The solution of this system is of the form

$$a_{1,a}^{(1)}(p)=D_{1,a}(p)/D(p),$$

$$D(p)=\left[ip-E_1-2\omega-\int dE\frac{|V_{1E}|^2}{ip-E-\omega}\right]\left[ip-E_a-\omega-\int dE\frac{|W_{aE}|^2}{ip-E-\omega}\right.$$
$$\left.-\int dE'\frac{|V_{aE'}|^2}{ip-E'}\right]-\left[V_{1a}^*+\int dE\frac{W_{aE}V_{1E}^*}{ip-E-\omega}\right]\left[V_{1a}+\int dE\frac{W_{aE}^*V_{1E}}{ip-E-\omega}\right]$$

$$D_1(p)=F_{0a}^*\left(V_{1a}+\int dE\frac{W_{aE}^*V_{1E}}{ip-E-\omega}\right)a_0^{(0)}(p),$$

$$D_a(p)=F_{0a}^*\left(ip-E_1-2\omega-\int dE\frac{|V_{1E}|^2}{ip-E-\omega}\right)a_0^{(0)}(p). \tag{6}$$

It is known that when the inverse Laplace transform is taken

$$a(t)=\frac{1}{2\pi i}\int_C a(p)e^{pt}dp \tag{7}$$

the integration contour C in the complex p plane, $p=p'+ip''$, passes to the right of all the singularities of the functions $a(p)$. An investigation of the analytic properties of the functions $a_1^{(1)}(p)$ and $a_a^{(1)}(p)$ can be carried out in perfect analogy with the investigation carried out in Ref. 7. As a result we find that these functions have no singularities in the entire complex p-plane, with the exception of a cut along the imaginary negative axis $p''<0$, $p'=0$. Using the known identity

$$\frac{1}{x\pm i0}=P\frac{1}{x}\mp i\pi\delta(x),$$

where P is the symbol of the principal value of the integral, we find that on the edges of the cut

$$D^{(\pm)}(p'')=\left(p''+\tilde{E}_1\mp\frac{i}{2}\Gamma_1\right)\left[p''+\tilde{E}_a\mp\frac{i}{2}(\Gamma_i+\Gamma_a)\right]-\left(\tilde{V}_{1a}^*\mp\frac{i}{2}\Gamma_{1a}^*\right)\left(\tilde{V}_{1a}\mp\frac{i}{2}\Gamma_{1a}\right), \tag{8}$$

$$D_1^{(\pm)}(p'')=\frac{iF_{0a}^*(\tilde{V}_{1a}\mp i\Gamma_{1a}/2)}{p''+\tilde{E}_0\mp i0},\quad D_a^{(\pm)}(p'')=\frac{iF_{0a}^*(p''+\tilde{E}_1\mp i\Gamma_1/2)}{p''+\tilde{E}_0\mp i0},$$

where the upper signs pertain to the right-hand ($p'\to+0$) and the lower to the left-hand edge of the cut. These equations enable us to find $a_1^{(1)\pm}(p'')$ and $a_a^{(1)\pm}(p'')$, and subsequently the function $a_E^{(\pm)}(p'')$ with the aid of Eq. (3). The following notation is used in the equations:

$$\tilde{E}_0=E_0+\Omega+\omega,\quad \tilde{E}_1=E_1+2\omega,\quad \tilde{E}_a=E_a+\omega,$$
$$\Gamma_1=2\pi|V_{1,-(p''+\omega)}|^2,\quad \Gamma_i=2\pi|V_{a,-p''}|^2,\quad \Gamma_a=2\pi|W_{a,-(p''+\omega)}|^2, \tag{9}$$

and the parameters Γ have the meaning of the level widths due to ionization of the states ψ_1 and of the AS in the strong field, and to the autoionization of the AS (Γ_a). We have left out of (8) the integrals, in the sense of the principal value, which determine the corrections to the energy values energy $E_{1,a}$. These integrals constitute part of the level shift due to the quadratic dynamic Stark effect, which we assume to be completely taken into account in the definitions of E_1 and E_a. Equations (8) include also "crossing" widths and the corrections δ_{1a} that must be introduced in the matrix element V_{1a} as a result of the indirect transitions:

$$\tilde{V}_{1a}=V_{1a}-\delta_{1a},\quad \delta_{1a}=\int dE\frac{V_{1E}W_{aE}^*}{p''+E+\omega},$$
$$\Gamma_{1a}=2\pi W_{aE}^*V_{1E}|_{E=-p''-\omega}. \tag{10}$$

All the produced widths and the quantity δ_{1a} are functions of the variable p''. This dependence, however, is slow (the characteristic scale over which a substantial change takes place in each of these functions in Ry). We shall therefore regard all these quantities as constants. This approximation is justified if the excess above the ionization threshold $\sim E_a$ is large compared with all the widths (see Ref. 16). In accordance with Eq. (3), the correction of first order in the weak field to the amplitude $a_0(p)$ is zero, and the second-order correction is given by

$$(ip-\tilde{E}_0)a_0^{(2)}(p)=F_{0a}a_a^{(1)}(p). \tag{11}$$

It is easy to show that the function $a_0^{(2)}(p)$ has the same analytic properties as the first-order functions discussed above. To find the function $a_0^{(2)}(t)$ it is necessary to calculate the integral (7), choosing an integration contour that passes to the right of the imaginary negative axis. We replace this integral by an integral along the imaginary negative axis $p''<0$, $p'=0$, of the discontinuity of the function $a_0^{(2)}(p)$ on the cut (see the analogous procedure in Ref. 7).

Without dwelling on the details of the calculations, we present the final expression for the atom ionization rate under the assumption

$$\operatorname{Im} x_{1,2}\cdot t\gg 1, \tag{12}$$

where $x_{1,2}$ are the roots of the equation $D^+(x)=0$, namely

$$x_{1,2}=-\frac{1}{2}\left[\tilde{E}_1-\frac{i}{2}\Gamma_1+\tilde{E}_a-\frac{i}{2}(\Gamma_a+\Gamma_i)\right]\pm\frac{1}{2}\left\{\left[\varepsilon+\frac{i}{2}(\Gamma_a+\Gamma_i-\Gamma_1)\right]^2+4\left(\tilde{V}_{1a}^*-\frac{i}{2}\Gamma_{1a}^*\right)\left(\tilde{V}_{1a}-\frac{i}{2}\Gamma_{1a}\right)\right\}^{1/2}. \tag{13}$$

Here $\varepsilon=E_1+\omega-E_a$ is the detuning from resonance in the system of levels E_1 and E_a.

Under the condition (12), the atom manages to go over during the pulse time from the states ψ_a and ψ_1 into the continuum. This means that all the atoms that leave the ground state ψ_0 are ionized. Therefore the atom excitation rate

$$\alpha=\frac{d}{dt}(1-|a_0(t)|^2)=-2\operatorname{Re}\dot{a}_0^{(2)}(t) \tag{14}$$

turns out to be equal to its ionization rate

$$\dot{w}=\int dE|a_E^{(1)}(t)|^2+\int dE'|a_{E'}^{(1)}(t)|^2. \tag{15}$$

The ionization probability depends in this case linearly on the time, $w=\alpha t$.

Assuming the matrix elements V_{1a} and Γ_{1a} to be real (in this case $\Gamma_{1a}^2=\Gamma_1\Gamma_a$), we obtain as a result of the calculation

$$\dot{w}=\alpha=|F_{0a}|^2\left[(\Gamma_a+\Gamma_i)\left(\varepsilon_1-\frac{1}{4}\frac{\Gamma_f\Gamma_{1a}}{\Gamma_a+\Gamma_i}\right)^2+\frac{1}{4}\left(\Gamma_1+\frac{1}{4}\frac{\Gamma_f^2}{\Gamma_i+\Gamma_a}\right)\times(\Gamma_1(\Gamma_i+\Gamma_a)-\Gamma_{1a}^2)\right]\left\{\left[\left(\varepsilon_1\varepsilon_a-\frac{1}{16}\Gamma_f^2\right)+\frac{1}{4}(\Gamma_{1a}^2-\Gamma_1(\Gamma_a+\Gamma_i))\right]^2+\frac{1}{4}\left[\varepsilon_1(\Gamma_a+\Gamma_i)+\varepsilon_a\Gamma_1-\frac{1}{2}\Gamma_f\Gamma_{1a}\right]^2\right\}^{-1}. \tag{16}$$

Here $\Gamma_f=4|\tilde{V}_{1a}|$ is the so called field width of the resonance curve (see Ref. 16), $\varepsilon_1=\tilde{E}_1-\tilde{E}_0+\omega-(E_0+\Omega)$ is the deviation from two-photon resonance in the system of levels E_1 and E_0, and $\varepsilon_a=\tilde{E}_a-\tilde{E}_0=E_a-(E_0+\Omega)$ is the deviation from resonance in the system of levels E_a and E_0.

As noted above, the expansion (1), and consequently Eq. (16), corresponds to the simplest possible situation, when the AS decays only into those continuum states which are coupled by the field V with the state E_1. Of course, cases are also possible when transitions from the AS (or from the state E_1) take place also into continuum states E that are not coupled with the state E_1 (or with the AS). It is easy to show by calculations similar to those described above that in these cases it is necessary to make in (16) the substitutions

$$\Gamma_a\to\bar{\Gamma}_a=\Gamma_a+\Gamma_a',\quad \Gamma_1\to\bar{\Gamma}_1=\Gamma_1+\Gamma_1'. \tag{17}$$

The parameters $\bar{\Gamma}_a$ and $\bar{\Gamma}_1$ are respectively the total autoionization width of the level E_a and the total ionization width of the level E_1; Γ_a and Γ_1 are those parts of these widths which are due to transitions to identical states of the continuum, i.e., the "interfering" parts of the autoionization and ionization widths; Γ_a' and Γ_1' are their "noninterfering" parts and are due to transitions from the state ψ_a into continuum states that are not coupled to ψ_1, or from the state ψ_1 into continuum states not coupled with ψ_a. The values of the parameters Γ_a' and Γ_1' are determined by the selection rules for the autoionization transition and for transitions in a strong field $\mathbf{E}$. In the general case Γ_a' and Γ_1' are not small compared with Γ_a and Γ_1, and these parameters must be taken into account (in contrast to the model formulation of the problem in Refs. 14 and 15). The allowance for the "noninterfering" widths Γ_a' and Γ_1' can be of fundamental significance, inasmuch as under certain conditions it is precisely these factors which can determine the minimum width of the autoionization resonance in a strong field (see Sec. 3 below). We note that when the substitutions (7) are made the expression for the crossing width Γ_{1a} remain unchanged: $\Gamma_{1a}^2 = \Gamma_1\Gamma_a$.

3. DEPENDENCE OF THE WIDTH OF THE RESONANCE CURVE ON THE INTENSITY AND ON THE FREQUENCY OF A STRONG FIELD

A. To investigate the behavior of the function α we consider first an experimental setup in which the frequencies ω and Ω of the two fields vary from measurement to measurement but their difference, i.e., the deviation from resonance in the system of levels E_1 and E_0, remains unchanged. In this case the absorption function α takes the form of a Lorentz curve as a function of the variable ε_a:

$$\alpha(\varepsilon_a)=|F_{0a}|^2\frac{\gamma(\varepsilon_1)}{(\varepsilon_a-\Delta(\varepsilon_1))^2+{}^1/_4\gamma^2(\varepsilon_1)},$$

$$\gamma(\varepsilon_1)=\left\{(\Gamma_i+\tilde{\Gamma}_a)\left(\varepsilon_1-\frac{1}{4}\frac{\Gamma_f\Gamma_{1a}}{\Gamma_i+\tilde{\Gamma}_a}\right)^2+\frac{1}{4}\left(\tilde{\Gamma}_1+\frac{1}{4}\frac{\Gamma_f^2}{\Gamma_i+\tilde{\Gamma}_a}\right)\cdot[\tilde{\Gamma}_1(\Gamma_i+\tilde{\Gamma}_a)-\Gamma_{1a}^2]\right\}\left(\varepsilon_1^2+\frac{1}{4}\tilde{\Gamma}_1^2\right)^{-1}, \tag{18}$$

$$\Delta(\varepsilon_1)=\frac{1}{4}\frac{\varepsilon_1({}^1/_4\Gamma_f^2-\Gamma_{1a}^2)+{}^1/_2\tilde{\Gamma}_1\Gamma_f\Gamma_{1a}}{\varepsilon_1^2+{}^1/_4\tilde{\Gamma}_1^2},$$

while the width γ and the position of the maximum $\varepsilon_a=\Delta$ of the resonance curve depend parametrically on the detuning ε_1 of the two-photon resonance. We now investigate in greater detail the parametric $\gamma(\varepsilon_1)$ dependence. As seen from (18), at large ε_1, $|\varepsilon_1|\gg\tilde{\Gamma}_1$ and $|\varepsilon_1|\gg\Gamma_f\tilde{\Gamma}_1/\Gamma_{1a}$, we have $\gamma(\varepsilon_1)\approx\Gamma_i+\tilde{\Gamma}_a$, i.e., in this case no effective population of the level E_1 takes place and the system decays mainly via the channel $E_0\to E_a\to(E, E')$. On the other hand, the minimum value of the width γ, which is reached at $\varepsilon_1=\Gamma_f\tilde{\Gamma}_1/4\Gamma_{1a}$, and the maximum ionization rate α, are given by

$$\gamma_{min}=\Gamma_i+\Gamma_a-\Gamma_{1a}^2/\Gamma_1,\quad \alpha_{max}=4|F_{0a}|^2/\gamma_{min}. \tag{19}$$

In this case $\Delta(\varepsilon_1)=\Gamma_f\Gamma_{1a}/4\Gamma_1$. If the field is strong enough, $\Gamma_1\sim\Gamma_a$, or if the widths Γ_1' and Γ_1' are not small, $\Gamma_a'\sim\Gamma_a$, $\Gamma_1'\sim\Gamma_1$, no substantial change of the resonance curve takes place, as can be seen from (19): $\gamma_{min}\sim\Gamma_a$. The resonance curve can be strongly narrowed in much weaker fields

$$\Gamma_a\gg\Gamma_i,\ \Gamma_1, \tag{20}$$

if, in addition, the "noninterfering" increments to the widths Γ_1 and Γ_a are small enough

$$\Gamma_a'\ll\Gamma_a,\quad \Gamma_1'\ll\Gamma_1. \tag{21}$$

Under these conditions $\gamma_{min}\approx\Gamma_i\ll\Gamma_a$. It can be stated that in this case the interference leads to a mutual extinction of the decay channels $E_0\to E_a\to E$ and $E_0\to E_a\to E_1\to E$, and the system decays effectively via the channel $E_0\to E_a\to E'$. This is accompanied also by a proportional growth in the maximum ionization rate (19) compared with the case when there is no strong field, when this rate is $4|F_{0a}|^2/\Gamma_a$. We note also that in this case the resonance in the level system E_1 and E_0 is much more exact at the maximum point than in the system E_a, E_0:

$$\varepsilon_1=\frac{1}{4}\Gamma_f\left(\frac{\Gamma_1}{\Gamma_a}\right)^{1/2}\ll\varepsilon_a=\Delta=\frac{1}{4}\Gamma_f\left(\frac{\Gamma_a}{\Gamma_1}\right)^{1/2}, \tag{22}$$

and it is this which ensures the possibility of effective interference of the two decay channels, satisfaction of the condition (20) notwithstanding.

B. Let now the experiment be organized in such a way that the frequency ω of the strong field is fixed. Using (16), it is convenient to represent the denominator of the function α in the form of a product of two quadratic functions:

$$\left[\left(\frac{\varepsilon_1+\varepsilon_a}{2}-\Delta(\varepsilon)\right)^2+\frac{(\gamma^+(\varepsilon))^2}{4}\right]\left[\left(\frac{\varepsilon_1+\varepsilon_a}{2}+\Delta(\varepsilon)\right)^2+\frac{(\gamma^-(\varepsilon))^2}{4}\right], \tag{23}$$

where

$$\Delta(\varepsilon)=\frac{1}{2}\operatorname{Re}\left(\left\{\left[\varepsilon+\frac{i}{2}(\tilde{\Gamma}_a+\Gamma_i-\tilde{\Gamma}_1)\right]^2+\left(\frac{1}{2}\Gamma_f-i\Gamma_{1a}\right)^2\right\}^{1/2}\right),$$

$$\gamma^\pm(\varepsilon)=\frac{1}{2}(\tilde{\Gamma}_a+\Gamma_i+\tilde{\Gamma}_1)\pm\frac{i}{2}\operatorname{Im}\left(\left\{\left[\varepsilon+\frac{i}{2}(\tilde{\Gamma}_a+\Gamma_i-\tilde{\Gamma}_1)\right]^2+\left(\frac{1}{2}\Gamma_f-i\Gamma_{1a}\right)^2\right\}^{1/2}\right), \tag{24}$$

and where Δ and $\gamma^\pm$ are the shifts and widths of the resonance maxima and depend parametrically on the detuning ε in the system of levels E_1, E_a or on the frequency ω of the strong field. The ionization rate α as a function of the detuning ε_a or of the frequency Ω of the weak field has two resonance maxima with widths $\gamma^\pm(\varepsilon)$, shifted by an amount of the order of Γ_f. It is easily seen that the width $\gamma^-(\varepsilon)$ has a minimum γ_{min} at $\varepsilon=\varepsilon_{min}$:

$$\varepsilon_{min}=-(\Gamma_f/4\Gamma_{1a})(\tilde{\Gamma}_a+\Gamma_i-\tilde{\Gamma}_1),$$
$$\gamma_{min}=\gamma^-(\varepsilon_{min})={}^1/_2\{\tilde{\Gamma}_1+\tilde{\Gamma}_a+\Gamma_i-[(\Gamma_i+\tilde{\Gamma}_a-\tilde{\Gamma}_1)^2+4\Gamma_{1a}^2]^{1/2}\}. \tag{25}$$

At the same value of the detuning ε_{min}, the width γ^+ of the second resonance maximum of the $\alpha(\Omega)$ curve reaches the maximum value γ_{max}:

$$\gamma_{max}=\gamma^+(\varepsilon_{min})={}^1/_2\{\Gamma_1+\Gamma_a+\Gamma_i+[(\Gamma_i+\Gamma_a-\Gamma_1)^2+4\Gamma_{1a}^2]^{1/2}\}.$$

At $\varepsilon=\varepsilon_{min}$, near the narrow maximum, the dependence of the absorption curve α on the frequency Ω of the weak field can be represented in the form:

$$\alpha(\varepsilon_a)=\frac{|F_{0a}|^2}{A^2}\left[\left(\varepsilon_1-\frac{1}{4}\frac{\Gamma_f\Gamma_{1a}}{\Gamma_i+\tilde{\Gamma}_a}\right)^2(\Gamma_i+\tilde{\Gamma}_a)+\frac{1}{4}\left(\tilde{\Gamma}_1+\frac{1}{4}\frac{\Gamma_f^2}{\Gamma_i+\tilde{\Gamma}_a}\right)\times\left(\tilde{\Gamma}_1(\Gamma_i+\tilde{\Gamma}_a)-\Gamma_{1a}^2\right)\right]\left\{\left[\varepsilon_a-\frac{A}{2}-\frac{1}{8}\left(\frac{\Gamma_f}{\Gamma_{1a}}(\Gamma_i+\tilde{\Gamma}_a-\tilde{\Gamma}_1)\right)\right]^2+\frac{\gamma_{min}^2}{4}\right\}^{-1}, \tag{26}$$

where

$$\varepsilon_1=\varepsilon_a-(\Gamma_f/4\Gamma_{1a})(\Gamma_i+\tilde{\Gamma}_a-\tilde{\Gamma}_1),$$
$$A=(\Gamma_f/4\Gamma_{1a})[(\Gamma_i+\tilde{\Gamma}_a-\tilde{\Gamma}_1)^2+4\Gamma_{1a}^2]^{1/2}.$$

We note that the minimum values of the widths γ_{min} [(19), (25)] are determined in both considered experimental setups by the ionization width Γ_i of the autoionizing state in the strong field, and by the "noninterfering" parts of the autoionization width Γ_a' and the ionization width Γ_1' of the discrete level. If all these factors are neglected, then $\gamma_{min}=0$, corresponding to the approximations and to the results of Refs. 14 (we recall that we do not take into account at all the radiative level widths). These assumptions, however, are not realistic. Whereas the parameters Γ_a' and Γ_1' can possibly be equal to zero for a certain special choice of the levels E_a, E_1, and E_0 and of the polarization of the strong field, in any case we always have in a strong field $\Gamma_i\neq 0$ and consequently $\gamma_{min}\neq 0$. As a rule, in this case γ_{min} is considerably larger than the characteristic scale of the radiative widths. Let the magnitude of the strong field be such that

$$\Gamma_1,\ \Gamma_i\sim\Gamma_a.$$

Then $\gamma_{min}\sim\Gamma_i\sim\tilde{\Gamma}_a$, i.e., in this case, just as in the analogous case of Sec. 3A, the interference of the channels does not lead to a substantial change of the width of the absorption curve $\alpha(\Omega)$.

A noticeable narrowing of the resonance curve is again possible in a weaker field

$$\Gamma_f\gg\Gamma_a\gg\Gamma_1,\ \Gamma_i \tag{27}$$

if the condition (21) is satisfied. [We note that the relations $\Gamma_f\gg\Gamma_i$ and $\Gamma_f\gg\Gamma_1$ are satisfied automatically, since $|\mathbf{E}_0|\ll E_{at}$ (Ref. 16).] Now, as seen from (25),

$$\gamma_{min}=\Gamma_1\Gamma_i/\Gamma_a\ll\Gamma_a, \tag{28}$$

i.e., one of the branches of the resonance curve can be greatly narrowed, and the maximum value α_{max} is again given by Eq. (19), while the detunings at $\alpha=\alpha_{max}$ satisfy the relations (22). Of course, the second branch of the $\alpha(\Omega)$ curve has the same width $\gamma^+(\varepsilon_{min})=\gamma_{max}\sim\Gamma_a$ as before, i.e., the maxima and the widths of both resonance branches differ greatly in this case.

We note here that the narrowing of the resonance curve and the increase of its amplitude can hardly be treated as a narrowing of the AS itself, i.e., as an increase of its lifetime. Under the indicated conditions, interference takes place between the transitions $E_0\to E_a\to E$ and $E_0\to E_a\to E_1\to E$ and leads to a narrowing of the absorption curve. The system raised to the state E_a, however, will decay, as can be easily seen, mainly via the channel $E_a\to E$ with a probability of the order of Γ_a per unit time, whereas the decay via the channel $E_a\to E_1\to E$ will have a much lower probability, which can be estimated from the equation for the resonant ionization probability[16] at

$$\frac{\Gamma_f^2\Gamma_1}{\varepsilon^2+{}^1/_4\Gamma_f^2}\approx\frac{\Gamma_f^2\Gamma_1}{\varepsilon^2}=\frac{\Gamma_1^2}{\Gamma_a}\ll\Gamma_a.$$

From the general expression (23) it follows that the widths $\gamma^{\pm}(\varepsilon)$ at large detunings $\varepsilon\gtrsim\Gamma_f\gg\tilde{\Gamma}_a,\ \tilde{\Gamma}_1,\ \Gamma_i$ are given by

$$\gamma^{\pm}(\varepsilon)=\frac{1}{2}\left\{\Gamma_a+\Gamma_i+\Gamma_1\pm\frac{[\varepsilon(\Gamma_a+\Gamma_i-\Gamma_1)-\Gamma_f\Gamma_{1a}]}{(\varepsilon^2+{}^1/_4\Gamma_f^2)^{1/2}}\right\},$$

from which it is seen that asymptotically we have

$$\gamma^{\pm}\to\begin{cases}\tilde{\Gamma}_1, & \varepsilon\to\mp\infty\\ \Gamma_i+\tilde{\Gamma}_a, & \varepsilon\to\pm\infty\end{cases},$$

i.e., at large deviation from resonance in the system E_a and E_1 the interference between the transitions $E_0\to E_a\to E$ and $E_0\to E_a\to E_1\to E$ vanishes and the decay via each channel is independent.

An approximate plot of the parametric dependence of the widths $\gamma^{\pm}$ on the detuning ε is shown in Fig. 2.

Thus, if we trace the modification of the absorption spectrum on the autoionization level as a function of the intensity of the resonant field, the qualitative relations consist of the following. In the absence of a strong field the absorption curve has one resonance maximum (dashed curve of Fig. 3). When a strong field is applied, the resonance is split; this corresponds to the usual Autler-Townes splitting in a two-level system E_a and E_1.[17,18] This splitting becomes noticeable at strong-field intensities such that $\Gamma_f\gtrsim\tilde{\Gamma}_a$. Generally speaking, the widths of both shifted maxima are this case of the same order as $\tilde{\Gamma}_a$. When the conditions (21) and (27) are satisfied, however, and when the resonances are specially tuned to the frequencies of the weak and strong fields, as determined by conditions (22), one of the maxima turns out to be anomalously narrowing with a width determined by Eq. (28) (Fig. 3). Hereafter, with increasing strong field, in the interval defined by the inequality (27), the narrow maximum broadens: its width increases up to a value $\sim\Gamma_a$, which is reached at Γ_i, $\Gamma_1\sim\Gamma_a$.

We consider now briefly the case when the weak-field frequency is such that $\hbar\Omega\approx E_1-E_0$, i.e., the atom becomes ionized in accordance with the scheme shown in Fig. 1(b).

Calculations similar to the preceding ones yield, when conditions (21) are satisfied, the following expressions for the ionization rate of the atom:

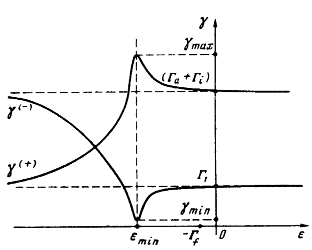

FIG. 2. Parametric dependence of the widths of the resonance curve of weak-field absorption energy $\alpha(\Omega)$ when the conditions (27) and (21) are satisfied.

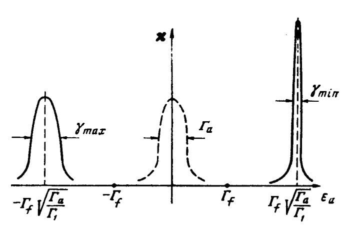

FIG. 3. Dependence of the absorption curve $\alpha(\Omega)$ on the weak-field frequency at $\varepsilon=\varepsilon_{min}=-1/4\Gamma_f(\Gamma_a/\Gamma_1)^{1/2}$ and $\Gamma_f\gg\Gamma_a\gg\Gamma_i,\ \Gamma_1,\ \Gamma_a'\ll\Gamma_a,\ \Gamma_1'\ll\Gamma_1$. The dashed curve shows the absorption in the absence of a strong field.

$$\dot{w}=|F_{01}|^2\frac{(\varepsilon_a\Gamma_1^{1/2}-{}^1/_4\Gamma_f\Gamma_a{}^{1/2})^2+{}^1/_4\,\Gamma_i[{}^1/_4\,\Gamma_f^2+\Gamma_1(\Gamma_i+\Gamma_a)]}{(\varepsilon_1\varepsilon_a-{}^1/_{16}\,\Gamma_f^2-{}^1/_4\,\Gamma_1\Gamma_i)^2+{}^1/_4[\varepsilon_1(\Gamma_a+\Gamma_i)+\varepsilon_a\Gamma_1-{}^1/_2\,\Gamma_f\Gamma_{1a}]^2}. \quad (29)$$

Here $\varepsilon_a=E_a-(E_0+\Omega+\omega)$, $\varepsilon_1=E_1-(E_0+\Omega)$, and the remaining symbols are the same as before.

Expression (29) is similar to Eq. (16), therefore all the foregoing results concerning the structure of the dispersion curves, the conditions for the narrowing of the resonance maxima, and the parameters that determine them remain in force also for resonant ionization of an atom in the presence of an additional resonance with AS.

4. CONCLUSION

Thus, in the considered interactions of AS with a discrete level (Fig. 1), the absorption spectrum of the probing field can be substantially altered by strong resonant radiation. If the frequencies are suitably chosen in the field interval defined by the inequalities (27), under the additional conditions (21) one of the resonance maxima can become substantially narrowed (see Fig. 3). Actual estimates of the corresponding intensities of the resonant field depend essentially on the value of the autoionization width Γ_a which varies in different atoms in a wide range. For example, at $\Gamma_a \sim 1\ \text{cm}^{-1}$ we have $\Gamma_f \sim \Gamma_a$ at strong-field intensity $|\mathbf{E}|\sim 5\cdot 10^4$ V/cm. It is clear from this estimate that the necessary fields are perfectly attainable and realizable, so that it can be concluded that the discussed regularities can be observed in experiment.

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