COHERENCE AND QUANTUM OPTICS VI

Proceedings of the Sixth Rochester Conference on Coherence and Quantum Optics held at the University of Rochester, June 26–28, 1989

Edited by

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Rochester, New York

PLENUM PRESS · NEW YORK AND LONDON

POPULATION TRAPPING EFFECT FOR SYSTEM WITH DOUBLE AUTOIONIZING LEVELS.

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INTRODUCTION

The existence of the population trapping effect for the ionization processes has been studied very extensively in many papers [1-6]. The special type of the ionization are the autoionization processes. They have been investigated in many papers (for example [7-13] and references quoted therein). One of the features associated with autoionization phenomena are effects corresponding to the existence of more than one of the autoionizing levels [14-18]. Those additional autoionizing levels and couplings between them can give various effects that do not occur for atomic systems containing single autoionizing level. For example, more than one zero can be visible in the photoelectron spectrum, additional peaks can occur, peaks can become broadened or be sharper (it depends on configuration of the system). Those additional zeros have different properties than the usual Fano-type zero [19]. They do not vanish even for large q-parameters. Moreover, the double confluence of coherences can be observable [16] for that kind of atomic system (the origin and nature of the confluence of coherences effect has been studied and explained by Rzazewski and Eberly [7]). Due to presence of more than one autoionizing level in the system the population trapping effect can appear too.

THE POPULATION TRAPPING EFFECT

We discuss an atomic system that contain two autoionizing levels $|1\rangle$ and $|2\rangle$ of the same parity (fig.1). Those levels are coupled to the ground discrete level $|0\rangle$ by the external laser field. We assume, that the laser field is monochromatic of frequency E_L (we use units of $h/2\pi$). Both autoionizing levels are diluted in the same continuum $|c\rangle$. Moreover, the laser field of the same frequency couples the ground state $|0\rangle$ to the continuum $|c\rangle$. All laser couplings are electric-dipole type only. Although the model discussed here is identical to that investigated in paper [16], our aim is to find conditions for the occurrence of the population trapping effect [1] rather, than to concentrate on the long-time photoelectron spectrum. Of course, the existence of the population trapping effect should manifest itself in the long-time photoelectron spectrum as a presence of sharp peaks.

Coherence and Quantum Optics VI
Edited by J. H. Eberly *et al.*
Plenum Press, New York, 1990

The atomic model discussed here is described by the following Hamiltonian (it is visible, that we do not perform Fano diagonalisation [16,19]):

$$H = |0\rangle(E_0 + E_L)\langle 0| + |1\rangle E_1\langle 1| + |2\rangle E_2\langle 2| + \int dE_c|c\rangle E_c\langle c| + \\ + \left(\int dE_c|c\rangle\Omega_c\langle 0| + \sum_{i=1,2}\left(|0\rangle\Omega_i\langle i| + \int dE_c|c\rangle V_i\langle i|\right. + \mathrm{H.c.}\right) \tag{1}$$

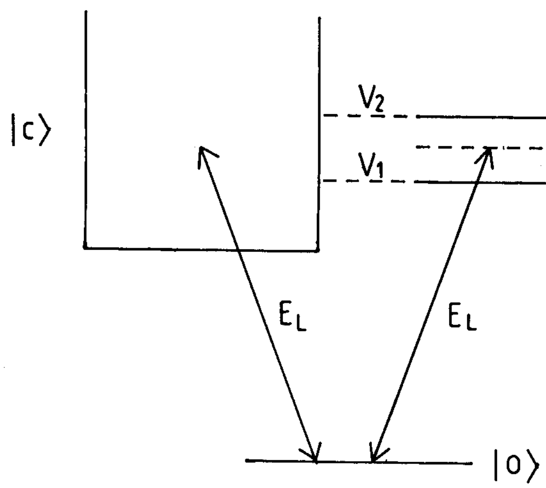

Fig. 1. The atomic model discussed in this paper.

where E_i (i=0,1,2) denotes energies of the levels $|0\rangle$, $|1\rangle$ and $|2\rangle$ respectively, E_c is energy of the continuum states $|c\rangle$, and V_i (i=1,2) corresponds to the configuration interaction between continuum $|c\rangle$ and autoionizing levels $|1\rangle$ and $|2\rangle$. The matrix elements (Rabi frequencies) Ω_c, Ω_i (i=1,2) characterize the laser field coupling between the continuum $|c\rangle$, the autoionizing levels $|1\rangle$ and $|2\rangle$ and to the ground state $|0\rangle$, respectively. For convenience we assume that all matrix elements are real. We neglect all threshold effects, so we extend lower limit of all integrals to minus infinity. It is equivalent to the assumption, that the both autoionizing levels $|1\rangle$ and $|2\rangle$ are far above the ionization threshold (for an example of the discussion of threshold effects see [20]). Moreover, we define the wave function $|\Phi\rangle$ for our system:

$$|\Phi(t)\rangle = a(t)\exp(-(E_o+E_L)t)|0\rangle + b_1(t)\exp(-E_1t)|1\rangle + \\ + b_2(t)\exp(-E_2t)|2\rangle + \int dE_c b_c(t)\exp(-E_ct)|c\rangle \tag{3}$$

We have introduced here the probability amplitudes a, b_1, b_2 and b_c. They are functions of time and correspond to the states $|0\rangle$, $|1\rangle$, $|2\rangle$ and $|c\rangle$ respectively. Moreover, one should keep in mind that the amplitude b_c is a function of the energy E_c too.

Performing the standard procedure based on the Schrodinger equation we can find the equations of motion (2) for the probability amplitudes a, b_i (i=1,2) and b_c. Moreover, we assume, that the laser has been turned on rapidly and that its amplitude remains constant during the process. The equations have the form:

$$i\frac{da}{dt} = \Omega_1 b_1 + \Omega_2 b_2 + \int dE_c \Omega_c b_c \tag{2a}$$

$$i\frac{db_1}{dt} = \delta_1 b_1 + \Omega_1 a + \int dE_c V_1 b_c \tag{2b}$$

$$i\frac{db_2}{dt} = \delta_2 b_2 + \Omega_2 a + \int dE_c V_2 b_c \tag{2c}$$

$$i\frac{db_c}{dt} = \Delta b_c + \Omega_c a + V_1 b_1 + V_2 b_2 \tag{2d}$$

where we have introduced the following detunings: $\delta_i = E_i - E_o - E_L$ (i=1,2), and $\Delta = E_c - E_O - E_L$. Solutions of the above equations (2a-2d) allow us to determine behaviour of the discussed system. Although it is possible to find the solutions for any strengths of the laser field, we restrict our considerations into the case of weak laser field couplings, i.e. we assume, that $\Omega_i, \Omega_c \ll V_i$ (i=1,2).

The first step is to eliminate b_c from the equations (2). Therefore, we find b_c from eqn.(2d) and put it into formulas (2a-2c). Assuming that matrix elements V_1, V_2 and Ω_c are smooth functions of the energy E_c (we see that matrix elements Ω_1 and Ω_2 do not depend on energy E_c), we can calculate all integrals appearing in equations (2) directly. Subsequently we can use the Laplace transform procedure to find the solutions of the above equations. Applying this method we obtain formulas that allow us to find the quasi-energies of the states $|0\rangle$, $|1\rangle$ and $|2\rangle$. One should keep in mind that these states are dressed as result of the interactions present in the system. The equations are of the form:

$$A(z)\{z + \Gamma_o\} + B_1(z)\{i\Omega_1 + \Gamma_{01}\} + B_2(z)\{i\Omega_2 + \Gamma_{02}\} = 1 \tag{3a}$$

$$A(z)\{i\Omega_1 + \Gamma_{01}\} + B_1(z)\{z + \Gamma_1\} + B(z)_2\Gamma_{12} = 0 \tag{3b}$$

$$A(z)\{i\Omega_2 + \Gamma_{02}\} + B_1(z)\Gamma_{12} + B_2(z)\{z + \Gamma_2\} = 0 \tag{3c}$$

where A(z), $B_1(z)$ and $B_2(z)$ are the Laplace transforms of the amplitudes a(t), $b_1(t)$ and $b_2(t)$, respectively. Moreover, we have defined here the following widths: $\Gamma_{Oi} = \pi\Omega_c V_i$, $\Gamma_i = \pi V_i^2$ (i=1,2), $\Gamma_{12} = \pi V_1 V_2$ and $\Gamma_o = \pi\Omega_c$. The widths Γ_i (i=1,2) are the autoionizing widths of the levels $|1\rangle$ and $|2\rangle$, respectively, Γ_o is the radiative width of the state $|0\rangle$, and Γ_{oi} (i=1,2) have both the radiative and autoionizing character. We have assumed that the discussed system was in its ground state $|0\rangle$ at the time $t=0$, i.e. $a(t=0)=1$ ($b_1=b_2=b_c=0$ for $t=0$). Although it is possible to find exact solutions of the equations (3a-3c), we will deal with the case of weak laser couplings. Therefore, we can apply the "pole approximation" [4] to find positions of the poles of the quantities $A(z)$, $B_1(z)$ and $B_2(z)$. It is very convenient to introduce at this point new complex matrix elements $\underline{\Omega}_1$ and $\underline{\Omega}_2$. They are defined as follows: $\underline{\Omega}_1 = i\Omega_1 + \Gamma_{01}$ and $\underline{\Omega}_2 = i\Omega_2 + \Gamma_{02}$. Applying the new matrix elements one can write down the zeros of the denominator of $A(z)$, $B_1(z)$ and $B_2(z)$. They are of the following form:

$$\epsilon_1 = \underline{\Omega}_1 \frac{\underline{\Omega}_1 - \underline{\Omega}_2 \dfrac{\underline{\Omega}_2 \Gamma_{12}}{i\delta_2 + \Gamma_2}}{i\delta_1 + \Gamma_1 - \dfrac{\Gamma_{12}^2}{i\delta_2 + \Gamma_2}} + \{1 \longleftrightarrow 2\} + \Gamma_o \qquad (4a)$$

$$\epsilon_{2,3} = -\frac{1}{2}\{i(\delta_1 + \delta_2) + \Gamma_1 + \Gamma_2\} \pm$$
$$\pm \{[i(\delta_1 + \delta_2) + \Gamma_1 - \Gamma_2]^2 + 4\Gamma_{12}^2\}^{1/2} \qquad (4b)$$

One of the quasi-energies ϵ_1 (eqn.4a) has rather complicated form that is not very interesting for our purposes.The remaining two of them - $\epsilon_{2,3}$ (eqn. 4b) are less complicated. They enable us to find the conditions for the occurrence of the population trapping effect. This effect should be visible when the quantity $(E_2-E_1)\Gamma_1\Gamma_2$ is equal to zero. One may see that the real part of one of the quasi-energies $\epsilon_{2,3}$ vanishes and population trapping should occur when the autoionizing width Γ_i ($i=1,2$) has zero value. One can explain this fact as a result of closing one of the autoionization channels. As a result the interference between radiative transitions can be essential. It is the same kind of the influence of the autoionizing width on the existing interferences between the transitions as investigated in the paper [18]). This interference gives the finite and nonzero value for the probability of finding the system in one of the dressed discrete levels, even for very long times. Therefore, one should expect a sharp peak in the long-time photoelectron spectrum that reflects the positions and widths of dressed atomic levels.

If energies E_1 and E_2 of the autoionizing levels $|1\rangle$ and $|2\rangle$ are the same, we may distinguish two cases: (i) that of the same parameters describing levels $|1\rangle$ and $|2\rangle$ ($\Gamma_1=\Gamma_2$, $\Omega_1=\Omega_2$) and (ii) that of autoionizing levels described by different autoionizing widths or (and) Rabi frequencies. For the case (i) our system reduces to that discussed by Rzazewski and Eberly [7], as it has been shown in [16]. When the autoionizing levels $|1\rangle$ and $|2\rangle$ are described by different parameters Γ_k and Ω_k ($k=1,2$) (ii), the population trapping effect should occur. One can suspect that for this situation the interference between transitions from the ground state $|0\rangle$ to autoionizing level $|1\rangle$ and from level $|0\rangle$ to $|2\rangle$ plays the most significant role.

ACKNOWLEDGMENTS

This work was supported by Research Project CPBP 01.07.

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