Laser-induced autoionizing and continuum structures: Line-shape study in the presence of continuum-continuum transitions

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(Received 9 September 1994)

We present a theoretical study of laser-induced autoionization and laser-induced continuum structures, performed with the aim of deriving predictions on the effect of continuum-continuum transitions on the resonance line shape. The model, developed as an extension of models introduced for the description of laser-induced autoionization and laser-induced continuum structures, on one hand, and for above-threshold ionization, on the other hand, provides quantitative predictions. For specific atomic systems, the line shapes of the structures are shown not to be affected by continuum-continuum transitions, in contrast with previously reported results.

PACS number(s): 32.80.Dz, 32.80.Fb, 32.70.Jz

I. INTRODUCTION

The study of the coherent atomic excitation to the ionization continuum has brought into evidence, in the photoionization process, some new effects caused by interference. It is well known, in fact, that, if different transition paths driving an atom to the same state in the continuum are allowed, quantum interference effects may drastically affect the transition rates and, under special conditions, the population can be trapped in a nondecaying state. For example, it has been pointed out recently that interference of transitions to the continuum from different Rydberg levels can be a reason for the stabilization of a highly excited atom: the time of ionization increases at increasing field intensity, up to the suppression of ionization, a rather unexpected result [1].

The first evidence of quantum interference in bound-continuum (B-C) transitions was provided by the observation of asymmetric peaks in continuous absorption spectra. The presence of autoionizing states was found responsible for these structures. A nonperturbative treatment of autoionization, suitable for a quantitative analysis of experimental results, was developed in 1961 by Fano [2]. The asymmetry of autoionizing resonances plays a crucial role in the atomic evolution, which may exhibit, for proper laser frequency and intensity, the phenomenon of population trapping [3,4]. In 1975, the possibility of inducing a structure in the photoionization spectrum, by a laser field coupling the continuum to a lower lying bound state, was predicted [5,6]. The bound state, embedded in the continuum by the intense radiation field (dressing field), behaves as a *pseudoautoionizing* state [5]: the compound *atom plus dressing-field* system, when probed by an additional weak radiation field (probe field), will exhibit a narrow structure in the photoionization spectrum, under many respects similar to an autoionizing resonance. This laser-induced continuum structure (LICS) was found very attractive for its potential in controlling the atomic photoionization rate by an external applied laser field, whose intensity and frequency can be arbitrarily varied. The analogy between LICS and laser-induced autoionization (LIA), proceeding from the observed effects to the physical interpretation, attains the formal description itself, which is almost unified in a theoretical model dealing with the excitation of structured continua [7]. This analogy is found to persist for the specific object of our study, namely, the effects of a strong radiation field inducing continuum-continuum (C-C) transitions on the line shape of the LICS and LIA resonances.

Experimental support to LICS theories has been pursued by several research groups undertaking experiments aimed at the measurement of various atomic properties. In fact, similarly to an autoionizing state, the bound state embedded in the continuum can affect, besides the photoionization rate, other atomic properties, such as, for instance, the birefringence, the dichroism, and the efficiency for harmonic generation, which can therefore be probed by specifically devoted experiments. The first experimental evidence of LICS was obtained in 1981 [8], through the measurement of the optical-polarization rotation induced by a σ-polarized dressing laser in cesium. Afterward, evidence of LICS was obtained by observing a resonant enhancement of third-harmonic generation in the continuum of sodium [9]. Considerable effort was also devoted to the direct observation of LICS in photoionization spectra, either in alkali metals [10] or in rare gases [11]. In 1991, a LICS resonance was observed in the single-photon ionization spectrum of sodium [12,13]. Evidence of LICS in two-electron atoms has been recently obtained studying an induced structure in the vicinity of an autoionizing state of calcium [14–16].

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In spite of the rich literature reporting LICS studies, further investigation both theoretical and experimental is required for a full physical understanding of the problem. In fact, the observation of LICS by third-harmonic generation in sodium, reported in [9], could not be reproduced by subsequent experiments [17,18], carried out on the same atom, but in somewhat different experimental conditions. Even the observation of LICS in photoionization spectra was not always successful, providing at times no evidence of the resonance [10] or results of difficult interpretation [11,19]. On the contrary, recent polarization experiments [20,21] in cesium and sodium, performed at a relatively low dressing-field intensity, have quantitatively confirmed the theoretical predictions. A reason for the unsatisfactory reproducibility of some experimental results can be found in the role played in a given experiment by concomitant processes, which can either raise the background or quench the resonance, thereby strongly reducing the accuracy of the measurement or even making the resonance not observable. It has been demonstrated, for instance, that the laser-induced enhancement of third-harmonic generation in the continuum depends, in actual experiment, on an interplay of various effects, including ionization saturation, that can drastically reduce the evidence of LICS [22]. Similarly, in multiphoton ionization experiments depending on the properties of the specific system, the background ionization can easily mask the LICS resonance. The high intensity of the dressing field required to induce the LICS in photoionization spectra, which can easily attain 10^{11} W/cm^2, makes the rejection of unwanted processes very difficult. Therefore a quantitative test of theories necessarily requires a careful knowledge of the atomic parameters influencing the LICS process, as well as a careful evaluation of background processes.

In concomitance to any among these processes, at field intensities typical of LICS experiments, above-threshold ionization (ATI) can take place [23]. A striking effect of *C-C* transitions on the total photoionization rate was predicted by Deng and Eberly in 1984 [24]. Considering the multiphoton ionization of a model atom, which was allowed to undergo *C-C* transitions, they predicted, by an analytically solvable model, a saturating behavior of the photoionization rate as a function of the laser intensity. It has been shown later that the strong-field dynamics of *C-C* transitions can be conveniently studied in two-color ATI experiments [25], where a low-frequency high-intensity laser (dressing laser) is used to couple continuum states, while a high-frequency low-intensity laser is used to ionize the atom. This configuration, allowing an independent variation of the *B-C* and *C-C* couplings, appears particularly well suited to the test of a possible saturation of the rate of photoionization towards multiple continua. A theoretical model describing a two-color ATI experiment has been proposed by Rzazewski, Wang, and Haus [26]: no saturation of the photoionization rate was shown by this model. These results, although showing that the atomic evolution can be critically dependent on the modeling of continuum interactions, have raised the question of how the line shape of continuum structures would be affected by *C-C* transitions. The problem has been recently discussed, both for LIA [27,28] and LICS [29]. Actually, both processes are expected to closely reflect the characteristics of the ATI dynamics. For instance, the resonant photoionization enhancement typical of LIA and LICS can exhibit a different dependence on laser intensity whether saturation of the photoionization rate is expected or not. Although, at present, the model development for these processes seems satisfactory, realistic predictions for actual atomic species and interaction conditions cannot be easily derived. The development of solvable models, as an extension of previous models developed for LIA and LICS, on one hand, and for ATI, on the other hand, is not straightforward. In fact, approximations that have been conveniently used in either model, such as, for instance, the flat-continuum approximation, used in many ATI treatments, or the rotating-wave approximation, used in LIA and LICS models, cannot be extended in general to a model describing both effects.

Our purpose here is to show that a solvable model for LICS and LIA, including *C-C* transitions, can be developed, with the advantage of providing predictions of line-shape effects for specific atomic systems and driving-field frequencies, while preserving the physical understanding of the problem. According to the conclusions of our survey, we will restrict our treatment to a single-photon probing of the continuum structure, modeling the processes on realistic assumptions for the atomic parameters.

In Sec. II we present the theoretical model, deriving the equations of motion which are then discussed for two limiting cases. In Sec. III the model is used to study the effect of *C-C* transitions on the LICS line shape. In this section a detailed discussion of the approximation introduced is also reported. Section IV is devoted to the study of the LIA line shape, in the presence of an additional laser field inducing *C-C* transitions. The conclusions are finally given in Sec. V.

II. THE LICS MODEL

We consider the photoionization of an atom with a structureless continuum, such as, for instance, an alkali atom not far from threshold. The atom is in the presence of two laser fields, one weak, of frequency ω_p (probe field), and the other strong, of frequency ω_d (dressing field) (Fig. 1). The probe field, typically in the UV spectral region, drives the atom, by a single-photon transition, from the ground state $|g\rangle$ to the continuum state $|1,\omega\rangle$, where 1 denotes a set of quantum numbers and ω the energy, measured from the ionization threshold. In the model, the multiphoton ionization induced by the dressing field, which is not expected to affect the line shape, is neglected. This process, of course, will give rise to a background ionization, which should be minimized for an accurate experimental study. The only relevant effect of the dressing field is assumed to be the coupling of excited bound states, initially unpopulated, to the ionization continuum. In fact, for a frequency difference $\omega_p-\omega_d$ corresponding to the energy difference between a given excited state $|e\rangle$ and the ground state $|g\rangle$, these states, assumed to have the same parity, are coupled by a two-photon transition, affecting the photoionization rate. The process can be described as the photoionization, induced by a weak probe field, of the atom *dressed* by a strong laser field, embedding the state $|e\rangle$ in the continuum $|1,\omega\rangle$. As a result of the $|g\rangle$-$|e\rangle$ coupling, a narrow feature will appear in the photoionization spectrum, as a function of the two-photon detuning $\Delta=[(\omega_p-\omega_d)-(\omega_e-\omega_g)]$. The $|g\rangle$-$|e\rangle$ transition can proceed either via intermediate continuum states, involving *B-C* radiative couplings, or via intermediate bound states $|k\rangle$, involving *B-B* couplings, as in usual discrete Raman processes. Therefore the whole atomic configuration will affect the properties of the photoionization structure, thereby making them unpredictably different from atom to atom. Discrete Raman processes will be explicitly included in our treatment, where they are shown to play a crucial role, according to previous works [30,31].

At increasing dressing-laser intensity, *C-C* transitions will take place. To take into account this effect, the model atom must include at least one additional ionization continuum $|2,\omega\rangle$, with a set of quantum numbers allowing dipole transitions $|1,\omega\rangle\leftrightarrow|2,\omega\rangle$, induced by the dressing field (Fig. 1). The overall process will then give rise to a multiple-peak structure of the electron energy spectrum, which might be somehow different from that of a pure ATI spectrum, in proximity to the two-photon resonance condition. Moreover, the occurrence of *C-C* transitions is expected to affect the LICS line shape. Both aspects of the problem have been the object of theoretical studies [29,32], leading to general predictions in weak- and strong-field regimes, mainly derived by numerical calculations. We will extend to LICS a solvable model which has been developed by Rzazewski, Wang, and Haus [26] for the treatment of the two-color ATI, where two separate lasers are used to ionize the atom and to induce transitions between continuum states. This configuration shows a close analogy to that occurring in LICS, where the dressing laser, in addition to redistributing the population among various continua, as it does in two-color ATI, couples one continuum to a bound state. We will show that predictions derived for the total ionization rate of the two-color ATI can be straightforwardly extended to LICS, providing information on the effect of *C-C* transitions on the resonance line shape.

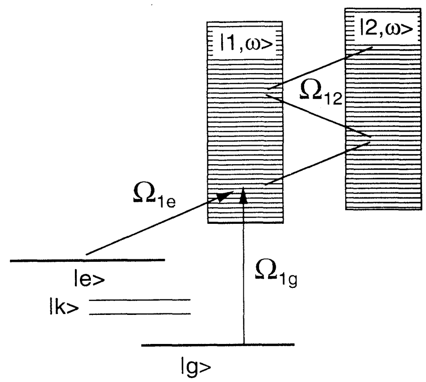

FIG. 1. Energy level scheme illustrating the LICS process in the presence of two manifolds of continuum states. Ω_{1e} is the Rabi frequency of a strong radiation field embedding the state $|e\rangle$ in the continuum $|1,\omega\rangle$. The same field couples continuum $|1,\omega\rangle$ to the continuum $|2,\omega\rangle$ by a Rabi frequency Ω_{12}. The continuum structure is probed by single-photon ground-state ionization, induced by a weak field of Rabi frequency Ω_{1g}.

The atomic state vector is expanded in terms of the bound and continuum states shown in Fig. 1 as

$$|\Psi(t)\rangle=c_g(t)|g\rangle+c_e(t)|e\rangle+\sum_k c_k(t)|k\rangle+\int d\omega c_1(t,\omega)|1,\omega\rangle+\int d\omega c_2(t,\omega)|2,\omega\rangle\ . \quad (1)$$

Following a semiclassical treatment of the atom-field interaction, the linearly polarized probe and dressing fields are written as

$$\mathbf{E}_p=E_p\mathbf{e}_p\cos(\omega_p t)\ \text{ and }\ \mathbf{E}_d=E_d\mathbf{e}_d\cos(\omega_d t)\ , \quad (2)$$

respectively, where $\mathbf{e}_p$ and $\mathbf{e}_d$ are unit polarization vectors. For the sake of an easy formalism, the field amplitudes E_p and E_d are assumed constant in time. By following the rotating-wave approximation (RWA) for the *B-C* transitions, the interaction Hamiltonian, in the Schrödinger picture, has the form

$$V=-e^{-i\omega_p t}\frac{\hbar}{2}\int d\omega\,\Omega_{1g}^{(p)}(\omega)|1,\omega\rangle\langle g|-e^{-i\omega_d t}\frac{\hbar}{2}\int d\omega\,\Omega_{1e}^{(d)}(\omega)|1,\omega\rangle\langle e|$$
$$-\cos(\omega_d t)\int d\omega\,d\omega'\hbar\Omega_{12}(\omega,\omega')|1,\omega\rangle\langle 2,\omega'|-\cos(\omega_p t)\sum_k \hbar(\Omega_{kg}^{(p)}|k\rangle\langle g|+\Omega_{ke}^{(p)}|k\rangle\langle e|)$$
$$-\cos(\omega_d t)\sum_k \hbar(\Omega_{kg}^{(d)}|k\rangle\langle g|+\Omega_{ke}^{(d)}|k\rangle\langle e|)+\text{H.c.} \quad (3)$$

where

$$\Omega_{kj}^{(p)}=\frac{1}{\hbar}E_p\langle k|\mathbf{d}|j\rangle\cdot\mathbf{e}_p\ \text{ and }\ \Omega_{kj}^{(d)}=\frac{1}{\hbar}E_d\langle k|\mathbf{d}|j\rangle\cdot\mathbf{e}_d\quad (j=g,e) \quad (4)$$

are *B-B* Rabi frequencies,

$$\Omega_{1g}^{(p)}=\frac{1}{\hbar}E_p\langle 1,\omega|\mathbf{d}|g\rangle\cdot\mathbf{e}_p\ \text{ and }\ \Omega_{1e}^{(d)}=\frac{1}{\hbar}E_d\langle 1,\omega|\mathbf{d}|e\rangle\cdot\mathbf{e}_d \quad (5)$$

are *B-C* Rabi frequencies, and

$$\Omega_{12}^{(d)}(\omega,\omega')=\frac{1}{\hbar}E_d\langle 1,\omega|\mathbf{d}|2,\omega'\rangle\cdot\mathbf{e}_d \tag{6}$$

is the *C-C* Rabi frequency. The upper script (p) or (d) assigned to a Rabi frequency denotes the laser field which is assumed to drive the corresponding transition.

The equations of motion for LICS are

$$\begin{aligned}
\dot c_g &= +\frac{i}{2}\int d\omega\,\Omega_{g1}^{(p)}(\omega)c_1(\omega)e^{-i(\omega-\omega_g-\omega_p)t}+i\sum_k[\Omega_{gk}^{(p)}\cos(\omega_p t)+\Omega_{gk}^{(d)}\cos(\omega_d t)]e^{i(\omega_g-\omega_k)t}c_k\ ,\\
\dot c_e &= \frac{i}{2}\int d\omega\,\Omega_{e1}^{(d)}(\omega)c_1(\omega)e^{-i(\omega-\omega_e-\omega_d)t}+i\sum_k[\Omega_{ek}^{(p)}\cos(\omega_p t)+\Omega_{ek}^{(d)}\cos(\omega_d t)]e^{i(\omega_e-\omega_k)t}c_k\ ,\\
\dot c_k &= i[\Omega_{kg}^{(p)}\cos(\omega_p t)+\Omega_{kg}^{(d)}\cos(\omega_d t)]e^{-i(\omega_g-\omega_k)t}c_g+i[\Omega_{ke}^{(p)}\cos(\omega_p t)+\Omega_{ke}^{(d)}\cos(\omega_d t)]e^{-i(\omega_e-\omega_k)t}c_e\ ,\\
\dot c_1(\omega) &= \frac{i\Omega_{1g}^{(p)}(\omega)}{2}e^{i(\omega-\omega_g-\omega_p)t}c_g+\frac{i\Omega_{1e}^{(d)}(\omega)}{2}c_e e^{i(\omega-\omega_e-\omega_d)t}+i\cos(\omega_d t)\int d\omega'\Omega_{12}^{(d)}(\omega,\omega')c_2(\omega')e^{i(\omega-\omega')t}\ ,\\
\dot c_2(\omega) &= i\cos(\omega_d t)\int d\omega'\Omega_{21}(\omega,\omega')c_1(\omega')e^{i(\omega-\omega')t}\ .
\end{aligned}\tag{7}$$

The dressing-laser frequency ω_d is required to be off resonance from the transitions $|g\rangle$-$|k\rangle$ and $|e\rangle$-$|k\rangle$, for any bound state $|k\rangle$, in order to avoid concomitant processes, which might be a source of background ionization. This requirement is obviously met by the probe laser frequency ω_p, which is large enough to induce single-photon ionization. The amplitudes c_k can be eliminated adiabatically from the equations of motion [33]. A further simplification can be introduced if we restrict the treatment to a small region around the two-photon resonance condition. In fact, with the phase transformation $c_{1,2}(\omega)\rightarrow c_{1,2}(\omega)e^{-i\omega t}; c_g\rightarrow c_g e^{-i(\omega_g+\omega_p)t}$ and the position $\omega_e+\omega_d=0$, for $\Delta=\omega_g+\omega_p\cong 0$, Eqs. (7) can be written as

$$\begin{aligned}
\dot c_g &= -i\Delta c_g+\frac{i}{2}\int d\omega\,\Omega_{g1}^{(p)}(\omega)c_1(\omega)+\frac{i}{2}\Omega_{ge}^{(R)}c_e\ ,\\
\dot c_e &= \frac{i}{2}\int d\omega\,\Omega_{e1}^{(d)}(\omega)c_1(\omega)+\frac{i}{2}\Omega_{eg}^{(R)}c_g\ ,\\
\dot c_1(\omega) &= -i\omega c_1(\omega)+\frac{i\Omega_{1g}^{(p)}(\omega)}{2}c_g+\frac{i\Omega_{1e}^{(d)}(\omega)}{2}c_e\\
&\quad +i\cos(\omega_d t)\int d\omega'\Omega_{12}^{(d)}(\omega,\omega')c_2(\omega')\ ,\\
\dot c_2(\omega) &= -i\omega c_2(\omega)+i\cos(\omega_d t)\int d\omega'\Omega_{21}^{(d)}(\omega,\omega')c_1(\omega')\ ,
\end{aligned}\tag{8}$$

where

$$\Omega_{ge}^{(R)}=\frac{1}{2}\sum_k\left[\frac{\Omega_{gk}^{(p)}\Omega_{ke}^{(d)}}{\omega_k-(\omega_e+\omega_d)}+\frac{\Omega_{gk}^{(d)}\Omega_{ke}^{(p)}}{\omega_k-(\omega_e-\omega_p)}\right]=\Omega_{eg}^{(R)*} \tag{9}$$

is the two-photon Rabi frequency corresponding to the Raman transition $|g\rangle-|e\rangle$. The Stark shift terms due to the states $|k\rangle$, not appearing explicitly in (8), can be assumed included in the detuning Δ. The energy zero has been taken at $\omega_e+\omega_d$.

Let us consider now two limiting cases.

Case (a): $\Omega_{1e}=0$. Equations (8) describe the two-color ATI, where the dressing field couples only the two continua between themselves. In this case, intermediate bound states can be safely neglected. Some simplifying assumptions make these equations exactly solvable [26]. The *B-C* coupling Ω_{g1} is assumed independent from the energy of continuum states. Moreover the *C-C* coupling $\Omega_{12}^{(d)}(\omega,\omega')$ is assumed to be a function of $\omega-\omega'$ alone. Finally, threshold effects are neglected. This assumption can be considered valid if the ionizing laser excites the electron well above threshold and the frequency of the dressing laser is small enough that transitions down to the threshold energy can be neglected. Retaining both frequency components of the dressing laser in *C-C* transitions has the effect of giving rise to a multipeak structure of the photoelectron spectra, even with only two continua, as is evident from the transition path shown in Fig. 1. But, what is most relevant to our problem, is the predicted behavior of the total photoionization rate versus the intensity of the dressing laser. Equations (8), with $\Omega_{1e}=0$ and $\Omega^{(R)}=0$, reduce to the equations of Ref. [26], where the photoionization rate turns out to be independent from the *C-C* coupling.

Case (b): $\Omega_{12}=0$. Equations (8) reduce to

$$\begin{aligned}
\dot c_g &= -i\Delta c_g+\frac{i}{2}\int d\omega\,\Omega_{g1}^{(p)}(\omega)c_1(\omega)+\frac{i}{2}\Omega_{ge}^{(R)}c_e\ ,\\
\dot c_e &= \frac{i}{2}\int d\omega\,\Omega_{e1}^{(d)}(\omega)c_1(\omega)+\frac{i}{2}\Omega_{eg}^{(R)}c_g\ ,\\
\dot c_1(\omega) &= -i\omega c_1(\omega)+\frac{i\Omega_{1g}^{(p)}(\omega)}{2}c_g+\frac{i\Omega_{1e}^{(d)}(\omega)}{2}c_e\ ,
\end{aligned}\tag{10}$$

describing the usual LICS, where the dressing laser embeds the state $|e\rangle$ in the continuum $|1,\omega\rangle$. The assumption of flat continuum, quite reasonable for ATI models, aimed at the study of the photoelectron spectrum, is particularly heavy for a LICS model, where the resonance profile itself is in general critically dependent on the energy behavior of *B-C* coupling factors. However, as will be shown below, for specific atomic species and laser frequencies, the characteristics of LICS can be negligibly affected by this energy dependence. Equations (10) can be solved approximately, for any energy dependence of Ω_{e1} and Ω_{g1}, provided that the continuum is widely extended to allow the use of a Markovian approximation in the B-C coupling [7]. In fact, under these assumptions, the continuum amplitudes can be eliminated, leading to the following equations for the bound-state amplitudes:

$$\dot{c}_g = \left[-i\Delta - \pi\frac{|\Omega_{1g}^{(p)}|^2}{4} \right] c_g - \frac{\pi}{4}\Omega_{g1}^{(p)}\Omega_{1e}^{(d)}(1-i\tilde{q})c_e \ ,$$
$$\dot{c}_e = -\pi\frac{|\Omega_{e1}^{(d)}|^2}{4}c_e - \frac{\pi}{4}\Omega_{e1}^{(d)}\Omega_{1g}^{(p)}(1-i\tilde{q})c_g \ , \tag{11}$$

where the shift terms have been again incorporated in Δ.

The quantity $\tilde{q}$, known as the Fano or asymmetry parameter, is an atomic parameter independent from the laser intensities. It can be written as $\tilde{q}=q+q^{(R)}$, with

$$q = \frac{1}{\pi\Omega_{g1}^{(p)}(\bar{\omega})\Omega_{1e}^{(d)}(\bar{\omega})} \mathrm{P}\int \frac{d\omega\Omega_{g1}^{(p)}(\omega)\Omega_{1e}^{(d)}(\omega)}{\omega-\bar{\omega}} \ , \tag{12}$$

$$q^{(R)} = \frac{2\Omega_{ge}^{(R)}}{\pi\Omega_{g1}^{(p)}(\bar{\omega})\Omega_{1e}^{(d)}(\bar{\omega})} \ , \tag{13}$$

and $\bar{\omega}\equiv\omega_e+\omega_d\cong\omega_g+\omega_p$. Equation (12) represents the contribution to the asymmetry parameter from the coupling to continuum states, while Eq. (13) represents the contribution from discrete Raman transitions. Notice that, in the flat-continuum approximation, owing to the vanishing of the principal part of the integral in (12), $q=0$ and consequently the line shape is determined by the configuration of bound states.

III. EFFECTS OF C-C TRANSITIONS ON THE LICS LINE SHAPE

We turn now back to the full set of equations (8). In order to make these equations solvable, we make the following assumptions.

(i) the B-C transitions are treated assuming both continua to be flat. With this assumption, implying $q=0$, the model can still be considered reliable for atomic systems meeting the condition $\tilde{q}\approx q^{(R)}$.

(ii) The C-C transitions are treated as in the two-color ATI model, assuming the corresponding Rabi frequency to be a real function, dependent only on the energy separation between the continuum states involved:

$$\Omega_{12}^{(d)}(\omega,\omega') = \Omega_{12}^{(d)}(|\omega-\omega'|) \ . \tag{14}$$

As a consequence of these assumptions threshold effects are neglected. Since in Eqs. (8) any transition is driven unambiguously by one laser field, we can safely set $\Omega_{g1}=\Omega_{g1}^{(p)}$, $\Omega_{e1}=\Omega_{e1}^{(d)}$, and $\Omega_{12}=\Omega_{12}^{(d)}$.

In order to take advantage of the energy dependence law (14), we introduce, following [26(b)], the Fourier transforms of the continuous energy variable:

$$\hat{c}_{1,2}(x,t) = \int_{-\infty}^{+\infty} d\omega\, e^{i\omega x} c_{1,2}(\omega,t) \ ,$$
$$\hat{\Omega}_{12}(x) = \int_{-\infty}^{+\infty} d\omega\, e^{i\omega x}\Omega_{12}(\omega) = \hat{\Omega}_{12}^{*}(x) \ . \tag{15}$$

The integro-differential Eqs. (8) are then transformed into a set of partial differential equations:

$$\dot{c}_g = -i\Delta c_g + \frac{i\Omega_{ge}^{(R)}}{2}c_e + \frac{i\Omega_{g1}}{2}\hat{c}_1(x=0,t) \ ,$$
$$\dot{c}_e = \frac{i\Omega_{eg}^{(R)}}{2}c_g + \frac{i\Omega_{e1}}{2}\hat{c}_1(x=0,t) \ ,$$
$$\frac{\partial\hat{c}_1(x,t)}{\partial t} = -\frac{\partial\hat{c}_1(x,t)}{\partial x} + \left[\frac{\Omega_{1g}}{2}c_g + \frac{\Omega_{1e}}{2}c_e \right] 2\pi i\delta(x) + i\hat{\Omega}_{12}(x)\hat{c}_2(x,t)\cos(\omega_d t) \ ,$$
$$\frac{\partial\hat{c}_2(x,t)}{\partial t} = -\frac{\partial\hat{c}_2(x,t)}{\partial x} + i\hat{\Omega}_{12}(x)\hat{c}_1(x,t)\cos(\omega_d t) \ , \tag{16}$$

which are solved according to Ref. [26(b)] providing, for the initial conditions

$$\hat{c}_1(t=0)=0, \quad \hat{c}_2(t=0)=0 \ , \tag{17}$$

the expressions

$$\hat{c}_1(x,t) = -2\pi i[\theta(x-t)-\theta(x)] \times \left[\frac{\Omega_{1g}}{2}c_g(t-x) + \frac{\Omega_{1e}}{2}c_e(t-x) \right] \times \cos\left\{ \int_0^x du\, \hat{\Omega}_{12}(u)\cos[\omega_d(u+t-x)] \right\} \ , \tag{18}$$

$$\hat{c}_2(x,t) = 2\pi[\theta(x-t)-\theta(x)] \times \left[\frac{\Omega_{1g}}{2}c_g(t-x) + \frac{\Omega_{1e}}{2}c_e(t-x) \right] \times \sin\left\{ \int_0^x du\, \hat{\Omega}_{12}(u)\cos[\omega_d(u+t-x)] \right\} \ ,$$

where $\theta(x)$ is the Heaviside unit step function.

Finally, by inserting these expressions in (16) we get, for the amplitudes of the discrete states, the equations

$$\dot{c}_g = -\left[i\Delta + \pi\frac{|\Omega_{g1}|^2}{4} \right] c_g - \frac{\pi}{4}\Omega_{g1}\Omega_{1e}(1-iq^{(R)})c_e \ ,$$
$$\dot{c}_e = -\pi\frac{|\Omega_{e1}|^2}{4}c_e - \frac{\pi}{4}\Omega_{e1}\Omega_{1g}(1-iq^{(R)})c_g \ , \tag{19}$$

provided that

$$\Omega_{12}(\omega-\omega') \xrightarrow[|\omega-\omega'|\to\infty]{} 0 \ ,$$

in order to avoid any divergence in its Fourier transform. These equations have the same form as the usual LICS equations (11) and do not contain the C-C coupling, implying that, under the assumptions of the model, the line shape of the resonance is not affected in any way by C-C transitions.

The effect of a flat-continuum assumption in the C-C coupling can be easily seen in Eqs. (18). In this simplifying but rather unphysical hypothesis, $\hat{\Omega}_{12}$ would be a δ function, $\hat{\Omega}_{12}(x)=2\pi\Omega_0\delta(x)$, and Eqs. (18) would lead, through a first-order development in Ω_0, valid for $\pi\Omega_0 \ll 1$, to the equations

$$\begin{aligned}\dot{c}_g &= -\left[i\Delta + \frac{\pi}{4}\frac{|\Omega_{g1}|^2}{1+\pi\Omega_0^2/4}\right]c_g \\ &\quad -\frac{\pi}{4}\frac{\Omega_{g1}\Omega_{1e}}{1+\pi\Omega_0^2/4}\left[1-iq^{(R)}\left[1+\frac{\pi\Omega_0^2}{4}\right]\right]c_e\ , \\ \dot{c}_e &= -\frac{\pi}{4}\frac{|\Omega_{e1}|^2}{1+\pi\Omega_0^2/4}c_e \\ &\quad -\frac{\pi}{4}\frac{\Omega_{e1}\Omega_{1g}}{1+\pi\Omega_0^2/4}\left[1-iq^{(R)}\left[1+\frac{\pi\Omega_0^2}{4}\right]\right]c_g\ .\end{aligned} \tag{20}$$

Equations (20) have been put in a form suitable to a direct comparison with (19). An effect of the *C-C* transitions on the LICS is now evident, both on the line shape, through a change of the asymmetry parameter, and on the photoionization rate.

It is worthwhile now to discuss the implications of some assumptions of the model, namely, (i) $\Omega_{12}^{(d)}(\omega,\omega')=\Omega_{12}^{(d)}(|\omega-\omega'|)$ and the consequent neglecting of threshold effects; (ii) inclusion of only two continua. Actually, the nonperturbative evaluation of the *C-C* coupling would require the inclusion of more than two continua. However, it has been shown that, even though the inclusion of more continua can affect the ATI peak amplitudes, it will not have any effect either on the LICS line shape or on the total ionization rate [34]. Assumption (i) seems more important since, as noted before, for

$$\Omega_{12}(\omega-\omega') \xrightarrow[|\omega-\omega'|\to\infty]{} 0\ ,$$

the model shows a LICS profile independent from *C-C* transitions. As a consequence any evidence of the dependence of the line shape on the *C-C* coupling would imply an effect of the ionization threshold.

Even the flat-continuum assumption in the treatment of *B-C* transitions, having different implications for the two-color ATI and for LICS, deserves a specific discussion. In ATI, only minor features of the photoelectron spectrum, which is the main object of study, are presumably affected by an energy dependence of *B-C* couplings, owing to the lack of any interference paths involving bound states. On the contrary, for LICS, the energy dependence of Ω_{g1} and Ω_{e1} is affecting the line shape itself, through the asymmetry parameter $\tilde{q}$. The appearance in Eqs. (19) of $q^{(R)}$ as the only contribution to the asymmetry parameter is just a consequence of the flat-continuum assumption in *B-C* transitions, necessary to make the equations solvable. Therefore this approximation can be considered valid for atomic systems and laser frequencies, for which $q^{(R)} \gg q$, so that $\tilde{q} \cong q^{(R)}$. Actually, numerical calculations [30] performed for cesium, for a configuration corresponding to an experimental study of LICS carried out through the measurement of the laser-induced polarization rotation [8,20], have shown that this condition is likely to occur. The bound states coupled to the continuum were $|g\rangle=(6s)^2S_{1/2}$ and $|e\rangle=(8s)^2S_{1/2}$, while the wavelengths of the dressing and probe fields were $\lambda_d=1064$ nm and $\lambda_p \simeq 296.6$ nm, respectively. The contribution to the Fano parameter from continuum states, q, turned out to be smaller than that from bound states, $q^{(R)}$, by more than one order of magnitude. In principle, the condition $q/q^{(R)} \ll 1$ might be reached for any atom by changing the position of induced resonances. In fact, due to the structure of the integrand, expression (12) can be minimized by a proper choice of the pole $\bar{\omega}$.

IV. EFFECTS OF *C-C* TRANSITIONS ON THE LIA LINE SHAPE

The model outlined above lends itself to the treatment of the laser-induced autoionization in the presence of a strong field inducing *C-C* transitions. The formal analogy between LIA and LICS will appear even closer for the physical system with which we are dealing. Referring to Fig. 2, we consider the transition paths typical of autoionization, with the addition of a second manifold of continuum states labeled $|2,\omega\rangle$ allowing, as in the case of LICS, *C-C* transitions. The coupling between the state $|e\rangle$ and the continuum states $|1,\omega\rangle$, denoted by T_{1e}, is due to an intra-atomic interaction instead of a dipole interaction. The atom is ionized by a weak probe field of frequency ω_p and Rabi frequency Ω_{g1}. Owing to the interference between different transition paths leading to the continuum $|1,\omega\rangle$, the autoionizing state gives rise to a resonance in the photoionization spectrum, characterized by an asymmetric profile, known as the Fano profile. We wish to study the effect, on the profile of the autoionizing resonance, of a strong field coupling the continua $|1,\omega\rangle$ and $|2,\omega\rangle$. Since the autoionizing state has the same parity as the states $|1,\omega\rangle$, the dressing field will also induce ionization of $|e\rangle$ towards the continuum $|2,\omega\rangle$. This process is physically different from the strong-field autoionization [27,28], where the coupling between different continua is induced by the same field which induces the ionization. In this respect, they can be said to

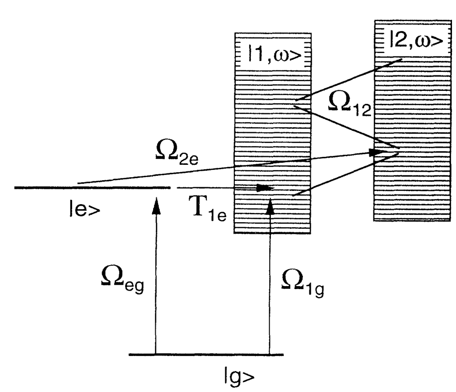

FIG. 2. Energy level scheme of the LIA process in the presence of two manifolds of continuum states. Here the state $|e\rangle$ is an autoionizing state coupled to the fundamental state $|g\rangle$ by the dipole interaction Ω_{eg} and to the continuum $|1,\omega\rangle$ by an intra-atomic interaction T_{1e}. The *C-C* transitions are induced by an additional field of Rabi frequency Ω_{12}. Since the autoionizing state is assumed to have the same parity as the states $|1,\omega\rangle$, the dressing field will also induce ionization of $|e\rangle$ towards the continuum $|2,\omega\rangle$.

be in the same relation, to each other, as the ATI and the two-color ATI. Actually, the LIA process in the presence of a dressing laser allows one to control the *B*-*C* and the *C*-*C* couplings separately, in analogy to the two-color ATI. Moreover, published works on strong-field autoionization have been mainly concerned with the study of the photoelectron spectrum, whereas our concern here is with the line shape of the autoionizing resonance in the total photoionization spectrum.

Since the effect of the bound states is now negligible, we can expand the state vector as

$$|\Psi(t)\rangle = c_g(t)|g\rangle + c_e(t)|e\rangle + \int d\omega\, c_1(t,\omega)|1,\omega\rangle + \int d\omega\, c_2(t,\omega)|2,\omega\rangle \ , \tag{21}$$

where $|e\rangle$ is the autoionizing state. By following the RWA for the $|g\rangle$-$|e\rangle$ and *B*-*C* transitions, the interaction Hamiltonian can be written as

$$V = -e^{-i\omega_p t}\frac{\hbar}{2}\int d\omega\, \Omega_{1g}^{(p)}(\omega)|1,\omega\rangle\langle g| + e^{-i\omega_p t}\frac{\hbar\Omega_{eg}}{2}|e\rangle\langle g| - \hbar\int d\omega\, T_{1e}|1,\omega\rangle\langle e| - \hbar\Omega_{e2}\cos(\omega_d t)\int d\omega|2,\omega\rangle\langle e| - \cos(\omega_d t)\int d\omega\, d\omega' \hbar\Omega_{12}(\omega,\omega')|1,\omega\rangle\langle 2,\omega'| + \text{H.c.} \ , \tag{22}$$

with

$$T_{e1} = \frac{1}{\hbar}\langle e|T|1,\omega\rangle \ . \tag{23}$$

The resulting equations of motion are

$$\begin{aligned}
\dot{c}_g &= -i\Delta c_g + \frac{i}{2}\int d\omega\, \Omega_{g1}(\omega)c_1(\omega) + i\frac{\Omega_{ge}}{2}c_e \ , \\
\dot{c}_e &= \frac{i}{2}\int d\omega\, T_{e1}c_1(\omega) + i\frac{\Omega_{eg}}{2}c_g + i\cos(\omega_d t)\int d\omega\, \Omega_{e2}(\omega)c_2(\omega) \ , \\
\dot{c}_1(\omega) &= -i\omega c_1(\omega) + \frac{i\Omega_{1g}(\omega)}{2}c_g + iT_{e1}c_e + i\cos(\omega_d t)\int d\omega' \Omega_{12}(\omega,\omega')c_2(\omega') \ , \\
\dot{c}_2(\omega) &= -i\omega c_2(\omega) + i\cos(\omega_d t)\int d\omega' \Omega_{21}(\omega,\omega')c_1(\omega') + i\Omega_{e2}(\omega)\cos(\omega_d t)c_e(\omega) \ .
\end{aligned} \tag{24}$$

The detuning is now defined as

$$\Delta = \omega_p - (\omega_e - \omega_g) \ . \tag{25}$$

Equations (24) can be treated by the same method and the same approximations used for LICS, i.e., assuming (i) Ω_{g1} and Ω_{e2} independent from the energy of continuum states, (ii) $\Omega(\omega,\omega') \equiv \Omega(|\omega-\omega'|)$. Notice that the flat-continuum approximation is supported by the presence of the direct $|g\rangle$-$|e\rangle$ coupling, implying [7]

$$\Omega_{ge} \gg P\int \frac{T_{e1}(\omega')\Omega_{g1}(\omega')}{\omega'-\omega_e}d\omega' \ . \tag{26}$$

By eliminating the amplitudes of the continuum states, we obtain

$$\begin{aligned}
\dot{c}_g &= -\left[i\Delta + \frac{\pi|\Omega_{g1}|^2}{4}\right]c_g + \left[\frac{i\Omega_{ge}}{2} - \frac{\Omega_{g1}}{2}T_{1e}\right]c_e \ , \\
\dot{c}_e &= -\pi[\,|T_{1e}|^2 + |\Omega_{e2}|^2\cos^2(\omega_d t)]c_e + \left[\frac{i\Omega_{eg}}{2} - T_{e1}\frac{\pi\Omega_{1g}}{2}\right]c_g \ .
\end{aligned} \tag{27}$$

These have the usual form of the LIA equations in the approximation of flat continuum, not containing the *C*-*C* coupling. The ionization of the state $|e\rangle$ to the continuum $|2,\omega\rangle$ gives rise to an additional time-dependent damping term, which averages to $|\Omega_{e2}|^2/2$. Therefore, even for the LIA process, the line shape of the resonance turns out to be independent from the *C*-*C* couplings.

V. CONCLUSION

We have described a study of the LIA and LICS processes, aimed at an evaluation of the effect of *C*-*C* transitions on the resonance line shape. For the LICS process, the *C*-*C* transitions are induced by the laser field embedding the bound state in the continuum, which in photoionization experiments can easily attain 10^{11} W/cm^2, whereas, for the LIA process, *C*-*C* transitions are assumed to be induced by an additional laser field coupling the manifolds of continuum states. We have shown that, restricting to specific atomic systems and well defined interaction schemes, quantitative predictions can be drawn by the model without numerical calculations. In particular, the physical system is required to present an asymmetry parameter mainly determined by coupling factors involving discrete states. The model shows that, under proper assumptions, for both processes, the total photoionization rate and the line shape of the resonance are independent from the intensity of the laser inducing *C*-*C* transitions, in contrast to previous results [29]. For the LIA process, the assumptions of the model are certainly fulfilled for autoionizing resonances, for which the dipole coupling between the ground and the autoionizing state is predominant over the coupling involving continuum states. For the LICS process, they can be met by systems where discrete Raman processes play a predominant role in the two-photon coupling.

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