# Strong-field effects in bichromatic laser-induced collisional energy transfer

Roberto Buffa, Stefano Cavalieri, Roberto Eramo, and Lorenzo Fini
*Dipartimento di Fisica and European Laboratory for Nonlinear Spectroscopy, Università di Firenze, Largo Enrico Fermi 2, 50125 Firenze, Italy*

Manlio Matera
*Istituto di Elettronica Quantistica, Consiglio Nazionale delle Ricerche, Via Panciatichi 56/30, 50127 Firenze, Italy*
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We discuss a scheme suitable for the study of strong-field effects on the excitation spectrum of the laser-induced collisional energy-transfer process. The use of a weak radiation field, to induce the interatomic transition, and a strong radiation field, to modify the energy-level position, is found to give rise to a splitting of the resonance peak. The position and the relative intensity of the two spectral components are strongly dependent on the intensity and detuning of the strong field. Line-shape calculations performed for the Eu-Sr system show that the effect is expected to be detectable even for moderate intensities of the strong field.

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## I. INTRODUCTION

Laser-induced collisional energy transfer (LICET) [1] is a process that involves the transfer of excitation from one atom to another with the simultaneous absorption of a photon. A typical configuration of atomic energy levels is shown in Fig. 1. Atom $A$, prepared in the excited state $|\alpha_2\rangle$, undergoes a collision with atom $B$, in its ground state $|\beta_1\rangle$, in the presence of a monochromatic laser field of frequency $\omega_p$ nearly resonant with the interatomic transition frequency $\omega_0=[E(|\beta_3\rangle)-E(|\alpha_2\rangle)]/\hbar$. Owing to the simultaneous action of the collisional and radiative interactions, one photon from the laser field is absorbed during the collision as the excitation energy of atom $A$ is transferred to atom $B$, which is then excited to state $|\beta_3\rangle$. The presence of an excited state $|\beta_2\rangle$ of atom $B$, nearly resonant with $|\alpha_2\rangle$, plays a crucial role in the collisional dynamics.

The earliest paper on LICET was probably the theoretical work reported by Gudzenko and Yakovlenko in 1972 [2], well before the first experimental observation by Falcone *et al.* in 1977 [3]. In the same year, Gallagher and Holstein [4] presented the first detailed calculations of LICET line shapes and the measurement reported by Brechignac, Cahuzac, and Toschek in 1980 [5] provided the first quantitative comparison between theory and experiments. Since then, theoretical models have undergone significant refinement [6,7], while the improvement of accuracy of further measurements [8–10] has allowed experiments to provide quantitative tests for the theory.

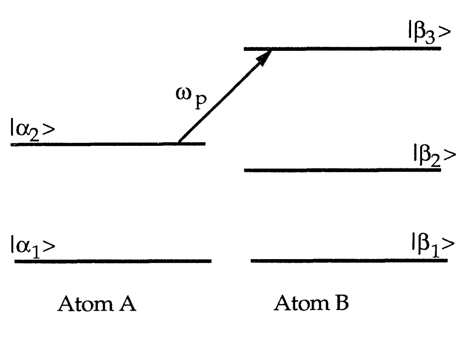


FIG. 1. Schematic energy-level diagram illustrating the LICET process.

For a weak laser intensity, the LICET cross section $\sigma(\omega_p)$, peaked at $\omega_p=\omega_0$, shows a strongly asymmetric shape with an extended wing related to the van der Waals shift of the atomic levels. The excellent agreement between theory and measurements recently obtained over the full line shape [10] has allowed a deep understanding of the process in the weak-field regime.

On the contrary, theoretical predictions of strong-field models [11] appear to lack experimental confirmation [12]. While at low intensity the radiation field acts just as a probe of adiabatic quasimolecular states formed during the collision, at increasing intensities it is expected to affect more deeply the dynamics of the process, dressing the atomic state $|\beta_2\rangle$ and $|\beta_3\rangle$. As a result, theoretical models predict a narrowing and symmetrization of the LICET line shape with a saturation and a frequency shift of the peak cross section [11]. However, so far, only the saturation of the peak cross section has been clearly confirmed by experiments [12].

These results raise the question whether the basic assumptions used in weak-field models are still valid at increasing laser intensities or the strong-field effects have not been observed because masked by concomitant effects, becoming increasingly important at increasing laser intensity. For instance, while it has been confirmed by experiments [10] that, in the weak-field regime, the LICET process can be correctly described by using an isotropic collisional potential, without being concerned with the magnetic degeneracy of states involved in the collision, it has been shown by Light and Szöke [13] that $m$-level degeneracy is overwhelmingly important for optical collisions in the strong-field regime. One could argue that, in the strong-field regime, the polarized laser radiation would dress the collisional molecule, distinguishing a preferred axis that would invalidate the assumption of spherically symmetric collision. However, strong-field LICET line shapes obtained in very recent calculations [14], performed including the sublevel degeneracy of states involved in the collision, confirm the frequency shift of the peak.

Moreover, at the laser intensity required to observe the frequency shift of the peak (larger than 100 MW/cm$^2$), physical processes, such as multiphoton transitions and ionization, or laboratory conditions, such as temporal and spatial laser pulse shape, can easily hide the predicted effect. For instance, following a simple rate-equation model, it can be argued that in a real experiment, performed in the Eu-Sr system with high vapor density ($10^{15}-10^{16}$ atoms/cm$^3$) and smooth laser pulses of 10–20 ns duration with intensity larger than 100–200 MW/cm$^2$, almost all the LICET reactions occur in the early rise time of the laser pulse and only a minor fraction of atoms experience the strong-field regime.

To overcome these problems, we have recently proposed a LICET experiment in a two-laser configuration [15]. In addition to the weak (probe) laser field of frequency $\omega_p$, a strong (dressing) laser field (not shown in Fig. 1) of frequency $\omega_d$, nearly resonant with the atomic transition $|\beta_2\rangle - |\beta_3\rangle$, is used to induce a dynamic Stark shift on these levels. The weak laser field will probe the collisional interaction involving atomic states dressed by the strong laser field. Moreover, since the dressing laser can be resonant with the transition $|\beta_2\rangle - |\beta_3\rangle$ or detuned by a small amount, a significant dynamic Stark effect can be induced by a moderate laser intensity, thereby reducing most of the problems of conventional strong-field LICET experiments.

In this paper we wish to discuss in more detail the bichromatic LICET process with the aim of predicting the dependence of the line shape on the dressing laser parameters (detuning and intensity). The different role played by the two laser fields allows a perturbative treatment of the equations of motion in a dressed-state basis, providing, under proper assumptions, explicit solutions. We will refer to the specific case of collisions between europium and strontium atoms, for which we have carried out an extensive investigation for the measurement of the weak-field LICET cross section [9,10].

## II. MODEL

A diagram of the relevant energy levels of the Eu and the Sr atoms involved in the bichromatic LICET process is shown in Fig. 2. The Eu atom, prepared in the $(6s6p)^8P_{9/2}$ excited state at 21 761 cm$^{-1}$, undergoes a collision with the Sr atom in the $(5s^2)^1S_0$ ground state in the presence of a weak laser field, nearly resonant with the interatomic transition $\mathrm{Eu}(6s6p)^8P_{9/2} \rightarrow \mathrm{Sr}(5p^2)^1D_2$, and of an intense laser field, nearly resonant with the Sr transition $(5s5p)^1P_1 \rightarrow (5p^2)^1D_2$. Since the $\mathrm{Sr}(5s5p)^1P_1$ excited state at 21 698 cm$^{-1}$ is nearly resonant with the $\mathrm{Eu}(6s6p)^8P_{9/2}$ initial state, one channel of excitation dominates the process, which is then conveniently studied in the following product-state basis:

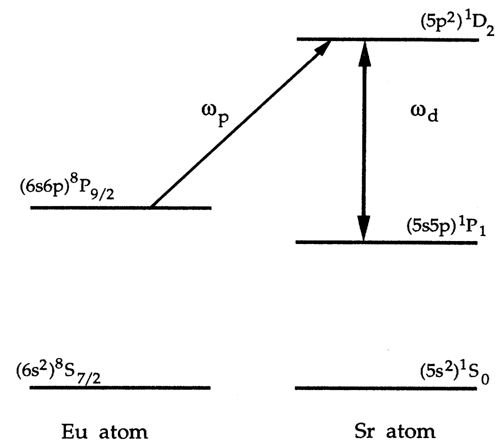


FIG. 2. Configuration of the relevant energy levels of the Eu and the Sr atoms involved in the bichromatic LICET process.

$$\begin{aligned} |1\rangle &= |\mathrm{Eu}(6s6p)^8P_{9/2}\rangle|\mathrm{Sr}(5s^2)^1S_0\rangle , \\ |2\rangle &= |\mathrm{Eu}(6s^2)^8S_{7/2}\rangle|\mathrm{Sr}(5s5p)^1P_1\rangle , \\ |3\rangle &= |\mathrm{Eu}(6s^2)^8S_{7/2}\rangle|\mathrm{Sr}(5p^2)^1D_2\rangle \end{aligned} \tag{2.1}$$

of energies $E_j = \hbar\omega_j$ ($j=1,2,3$). Following the approach reviewed in Ref. [1], several assumptions are made: (i) the collisional coupling is assumed to be a long-range electrostatic interaction, (ii) the trajectories for the colliding atoms are assumed classical and rectilinear, (iii) the electromagnetic field is assumed to be classical, and (iv) the magnetic degeneracy of the atomic states involved in the process is ignored. The equations of motion for the probability amplitudes of states (2.1) are then written as

$$i\dot{\mathbf{a}} = \underline{A}\,\mathbf{a} \tag{2.2}$$

with

$$\mathbf{a} = \begin{bmatrix} a_1 \\ a_2 \\ a_3 \end{bmatrix} \tag{2.3}$$

and

$$\underline{A} = \begin{bmatrix} \omega_1 & V & 0 \\ V & \omega_2 & -\chi \\ 0 & -\chi^* & \omega_3 \end{bmatrix} . \tag{2.4}$$

The dipole-dipole collisional interaction $V$ is given by

$$V = d_{\mathrm{Eu}} d_{\mathrm{Sr}} / \hbar (b^2 + v^2 t^2)^{3/2} , \tag{2.5}$$

where $b$ is the impact parameter, $v$ is the relative speed of the atoms, and $d_{\rm Eu}$ and $d_{\rm Sr}$ are the electrical-dipole moments of the Eu $(6s^2)^8S_{7/2} \rightarrow (6s6p)^8P_{9/2}$ and the Sr $(5s^2)^1S_0 \rightarrow (5s5p)^1P_1$ transitions, respectively. The electric-dipole radiative interaction $\chi$ is written, in the rotating-wave approximation, as

$$\chi = \chi_d \exp(i\omega_d t) + \chi_p \exp(i\omega_p t) , \tag{2.6}$$

where $\chi_d$ and $\chi_p$ are the Rabi frequencies of the dressing and the probe laser fields, respectively. As a result of the short time scale of the collisional interaction, of the order of 1–10 ps, $\chi_d$ and $\chi_p$ can be assumed constant during the single collision.

The coefficients of the wave function expansion (2.3) will have a rapidly varying time-dependent phase. However, it has been shown [6] that, when $\dot{V}/V \ll \omega_1 - \omega_2$ (63 $\rm cm^{-1}$ for Eu and Sr), this has no effect on the absorption spectrum of the process since no real transfer of population among states (2.1) can be induced by the collision alone (adiabatic interaction).

The cross section of the process is defined as

$$\sigma(\omega_p) = \left\langle \int_0^{+\infty} 2\pi b [1 - |a_1(+\infty)|^2] db \right\rangle_v , \tag{2.7}$$

where averages over the impact parameter $b$ and relative speed $v$ are carried out. Equation (2.2) can be numerically integrated starting from the initial conditions

$$\mathbf{a}(-\infty) = \begin{bmatrix} 1 \\ 0 \\ 0 \end{bmatrix} , \tag{2.8}$$

providing the solution $a_1(+\infty)$, which is then introduced in (2.7) to calculate the cross section. However, the different role played by the two laser fields allows a perturbative treatment of the equations of motion in a radiative dressed-state basis [16] defined by

$$\begin{aligned} |2d\rangle &= \cos\theta|2\rangle - \sin\theta \exp(-i\omega_d t)|3\rangle , \\ |3d\rangle &= \sin\theta|2\rangle + \cos\theta \exp(-i\omega_d t)|3\rangle , \end{aligned} \tag{2.9}$$

with $\tan(2\theta) = 2\chi_d/\Delta$ and $\Delta = \omega_d + \omega_2 - \omega_3$.

The approach of dressing the collisional molecule is not new. It had been introduced independently by Gudzenko and Yakovlenko [2] and by Kroll and Watson [17] and it was used by Light and Szöke to study the effect of $m$-level degeneracy in strong-field optical collisions [13].

Then, in terms of the new variables $\mathbf{X}$, defined by the unitary transformation

$$\mathbf{X} = \begin{bmatrix} 1 & 0 & 0 \\ 0 & \cos\theta & -\sin\theta \\ 0 & \sin\theta & \cos\theta \end{bmatrix} \begin{bmatrix} a_1 \\ a_2 \\ a_3 \exp(i\omega_d t) \end{bmatrix} , \tag{2.10}$$

the equations of motion (2.2) read as

$$i\dot{\mathbf{X}} = \underline{H}\mathbf{X} \tag{2.11}$$

with

$$\underline{H} = \begin{bmatrix} \omega_1 & V\cos\theta & V\sin\theta \\ V\cos\theta & \lambda_2 + H_{22}^{(1)} & H_{23}^{(1)} \\ V\sin\theta & H_{32}^{(1)} & \lambda_3 + H_{33}^{(1)} \end{bmatrix} \tag{2.12}$$

and where

$$\begin{aligned} \lambda_2 &= (\cos\theta)^2\omega_2 + \sin(2\theta)\chi_d + (\sin\theta)^2(\omega_3 - \omega_d) , \\ \lambda_3 &= (\sin\theta)^2\omega_2 - \sin(2\theta)\chi_d + (\cos\theta)^2(\omega_3 - \omega_d) , \\ H_{22}^{(1)} &= -H_{33}^{(1)} = \chi_p \sin\theta\cos\theta\{\exp[i(\omega_d - \omega_p)t] \\ &\qquad + \exp[i(\omega_p - \omega_d)t]\} , \\ H_{23}^{(1)} &= H_{32}^{(1)*} = \chi_p\{(\sin\theta)^2\exp[i(\omega_d - \omega_p)t] \\ &\qquad - (\cos\theta)^2\exp[i(\omega_p - \omega_d)t]\} . \end{aligned} \tag{2.13}$$

In the spirit of perturbation theory, we write

$$\underline{H} = \underline{H}^{(0)} + \underline{H}^{(1)} \tag{2.14}$$

with

$$\underline{H}^{(0)} = \begin{bmatrix} \omega_1 & V\cos\theta & V\sin\theta \\ V\cos\theta & \lambda_2 & 0 \\ V\sin\theta & 0 & \lambda_3 \end{bmatrix} , \tag{2.15}$$

and

$$\underline{H}^{(1)} = \begin{bmatrix} 0 & 0 & 0 \\ 0 & H_{22}^{(1)} & H_{23}^{(1)} \\ 0 & H_{32}^{(1)} & H_{33}^{(1)} \end{bmatrix} , \tag{2.16}$$

treating $\underline{H}^{(1)}$ as a weak perturbation. Thus, writing $\mathbf{X} = \mathbf{X}^{(0)} + \mathbf{X}^{(1)}$, we obtain, to the first order in $\chi_p$,

$$i\dot{\mathbf{X}}^{(0)} = \underline{H}^{(0)}\mathbf{X}^{(0)} \tag{2.17}$$

$$i\dot{\mathbf{X}}^{(1)} = \underline{H}^{(0)}\mathbf{X}^{(1)} + \underline{H}^{(1)}\mathbf{X}^{(0)} \tag{2.18}$$

with boundary conditions given by

$$\mathbf{X}^{(0)}(-\infty) = \begin{bmatrix} 1 \\ 0 \\ 0 \end{bmatrix}, \quad \mathbf{X}^{(1)}(-\infty) = \begin{bmatrix} 0 \\ 0 \\ 0 \end{bmatrix} . \tag{2.19}$$

The solution of (2.18) can be written as

$$\mathbf{X}^{(1)} = i\underline{U}\int_{-\infty}^{t} \underline{U}^{-1}\underline{H}^{(1)}\mathbf{X}^{(0)}dt' , \tag{2.20}$$

where $\underline{U}$ is the time-evolution operator, the solution of

$$i\dot{\underline{U}} = \underline{H}_0\underline{U} . \tag{2.21}$$

Since the dressing laser field is strongly detuned to the antistatic side of the LICET excitation spectrum, no real transfer of population can be induced by the strong laser alone, even though states (2.9) can be populated during the collision. As a result

$$\mathbf{X}^{(0)}(+\infty) = \begin{bmatrix} 1 \\ 0 \\ 0 \end{bmatrix} \tag{2.22}$$

and

$$\sigma(\omega_p)=\left\langle \int_0^{+\infty} 2\pi b[|X_2^{(1)}(+\infty)|^2+|X_3^{(1)}(+\infty)|^2]db \right\rangle_v . \quad (2.23)$$

## III. ADIABATIC APPROXIMATION

The Hamiltonian $\underline{H}_0$ can be diagonalized by the unitary transformation

$$\underline{D}\,\underline{H}_0\underline{D}^{-1}=\underline{M}=\begin{bmatrix}\mu_1 & 0 & 0\\ 0 & \mu_2 & 0\\ 0 & 0 & \mu_3\end{bmatrix} , \quad (3.1)$$

where the matrix elements $D_{ij}$ of $\underline{D}$ are given by

$$D_{i1}=(\mu_i-\lambda_2)(\mu_i-\lambda_3)/S_i ,$$
$$D_{i2}=\cos\theta V(\mu_1-\lambda_3)/S_i , \quad (3.2)$$
$$D_{i3}=\sin\theta V(\mu_i-\lambda_2)/S_i$$

with

$$S_i^2=[(\mu_i-\lambda_2)(\mu_i-\lambda_3)]^2+[\cos\theta V(\mu_i-\lambda_3)]^2 +[\sin\theta V(\mu_i-\lambda_2)]^2 \quad (3.3)$$

and where the time-dependent eigenvalues $\mu_i$ are given by the solution of the third-order secular equation

$$|\underline{H}_0-\mu\underline{I}|=0 . \quad (3.4)$$

Diagonalization (3.1) leads to the definition of the following radiative-collisional adiabatic dressed states:

$$|D_i\rangle=D_{i1}|1\rangle+D_{i2}|2d\rangle+D_{i3}|3d\rangle . \quad (3.5)$$

For $V=0$,

$$\mu_1=\omega_1,\quad \mu_2=\lambda_2,\quad \mu_3=\lambda_3 \quad (3.6)$$

and

$$|D_1\rangle=|1\rangle,\quad |D_2\rangle=|2d\rangle,\quad |D_3\rangle=|3d\rangle . \quad (3.7)$$

Therefore, in the absence of collision, state $|D_2\rangle$ lies above (below) state $|D_3\rangle$ for $\Delta>0$ ($\Delta<0$). During the collision, owing to the van der Waals shift of levels $|1\rangle$ and $|2\rangle$, state $|D_1\rangle$ is pushed upward, while states $|D_2\rangle$ and $|D_3\rangle$ are both pushed downward. However, an analysis of the temporal evolution of eigenstates $\mu_i$ during the collision shows that, while for $\Delta>0$, $\mu_2$ and $\mu_3$ present an avoided crossing, for $\Delta<0$ the eigenstates $\mu_i$ never cross, suggesting that an adiabatic analysis of the interaction dynamics can be appropriate for negative values of $\Delta$.

The matrix $\underline{DU}$ evolves in time according to

$$i(\dot{\underline{DU}})=(\underline{M}+i\dot{\underline{D}}\underline{D}^{-1})(\underline{DU}) . \quad (3.8)$$

If the generalized Rabi frequency $\Omega$ associated with the radiative interaction of the strong laser with the Sr atom ($\Omega=\sqrt{\Delta^2+4\chi_d^2}=|\lambda_2-\lambda_3|$) is sufficiently larger than the width $\Delta\omega$ of the Fourier transform of the collisional interaction ($\Delta\omega\approx v/2\pi b$), then the off-diagonal terms of $\dot{\underline{D}}\underline{D}^{-1}$ in (3.8) are negligibly smaller than $|\mu_i-\mu_j|$ ($i\neq j$) at any time. We can therefore ignore the term $\dot{\underline{D}}\underline{D}^{-1}$ in (3.8) to find the following solution for $\underline{U}$:

$$\underline{U}=[U_{ij}]=\left[D_{ji}\exp\left[-i\int_{-\infty}^{t}\mu_j dt'\right]\right] . \quad (3.9)$$

Since for $t=-\infty$, $\underline{U}=\underline{I}$, the first column of (3.9) is the solution of (2.17) to be introduced in (2.20) and the following solutions for $|X_i^{(1)}(+\infty)|^2$ are then obtained:

$$|X_2^{(1)}(+\infty)|^2=\chi_p^2\left|\int_{-\infty}^{+\infty}dt[(\cos\theta)D_{12}+(\sin\theta)D_{13}][(\sin\theta)D_{22}-(\cos\theta)D_{23}]\exp\left[i\int_{-\infty}^{t}(\mu_2+\omega_d)-(\mu_1+\omega_p)dt'\right]\right|^2$$
$$=\chi_p^2\chi_d^2\left|\int_{-\infty}^{+\infty}dt\, V^2\frac{(\mu_1-\omega_2+\Delta)}{S_1S_2}\exp\left[i\int_{-\infty}^{t}(\mu_2+\omega_d)-(\mu_1+\omega_p)dt'\right]\right|^2 , \quad (3.10)$$

$$|X_3^{(1)}(+\infty)|^2=\chi_p^2\left|\int_{-\infty}^{+\infty}dt[(\cos\theta)D_{12}+(\sin\theta)D_{13}][(\sin\theta)D_{32}-(\cos\theta)D_{33}]\exp\left[i\int_{-\infty}^{t}(\mu_3+\omega_d)-(\mu_1+\omega_p)dt'\right]\right|^2$$
$$=\chi_p^2\chi_d^2\left|\int_{-\infty}^{+\infty}dt\, V^2\frac{(\mu_1-\omega_2+\Delta)}{S_1S_3}\exp\left[i\int_{-\infty}^{t}(\mu_3+\omega_d)-(\mu_1+\omega_p)dt'\right]\right|^2 .$$

The bichromatic LICET excitation spectrum is then expected to exhibit a two-peak structure, with lines centered at the frequencies

$$\omega_p=\omega_d+\lambda_2-\omega_1,\quad \omega_p=\omega_d+\lambda_3-\omega_1 \quad (3.11)$$

and separated by the generalized Rabi frequency $\Omega=\sqrt{\Delta^2+4\chi_d^2}=|\lambda_2-\lambda_3|$. Near the peaks the main contribution to the cross section is due to collisions with large impact parameters, so that a perturbative development in the collisional interaction is reasonable in this region and, to the first order, the following expressions for (3.10) are obtained:

$$|X_2^{(1)}(+\infty)|^2 = \chi_p^2(\sin\theta)^2\left[\frac{\Delta+\omega_{12}}{\omega_{12}(\omega_{12}+\Delta)-\chi_d^2}\right]^2 \times\left|\int_{-\infty}^{+\infty} V\exp[i(\omega_d+\lambda_2-\omega_p-\omega_1)t]dt\right|^2, \tag{3.12}$$

$$|X_3^{(1)}(+\infty)|^2 = \chi_p^2(\cos\theta)^2\left[\frac{\Delta+\omega_{12}}{\omega_{12}(\omega_{12}+\Delta)-\chi_d^2}\right]^2 \times\left|\int_{-\infty}^{+\infty} V\exp[i(\omega_d+\lambda_3-\omega_p-\omega_1)t]dt\right|^2.$$

Therefore, for detunings $\Delta$ and Rabi frequencies $\chi_d$ not too large compared with the energy defect $\omega_{12}$ $(=63\ \text{cm}^{-1})$, we expect that the intensity of the peak centered at $\omega_p=\omega_d+\lambda_2-\omega_1$ $(\omega_p=\omega_d+\lambda_3-\omega_1)$ will depend on the dressing laser parameters as $\sin^2\theta(\cos^2\theta)$ does: increasing (decreasing) with the Rabi frequency $\chi_d$ and decreasing (increasing) with detuning $\Delta$.

## IV. RESULTS AND DISCUSSION

Expressions (3.3) have been used in (3.10) to calculate the excitation spectrum (2.23). All calculations reported here have been with a collisional interaction energy $d_{\text{Eu}}d_{\text{Sr}}=2.17\times10^{-35}$ erg cm$^3$ and taking the relative

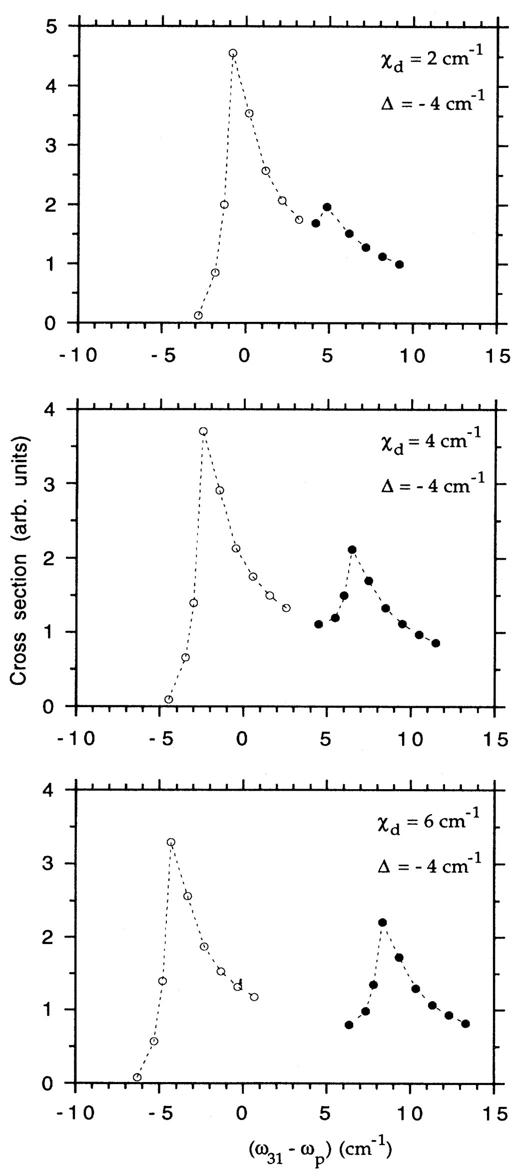


FIG. 3. Dependence of the bichromatic LICET line shape on strong-field Rabi frequency $\chi_d$.

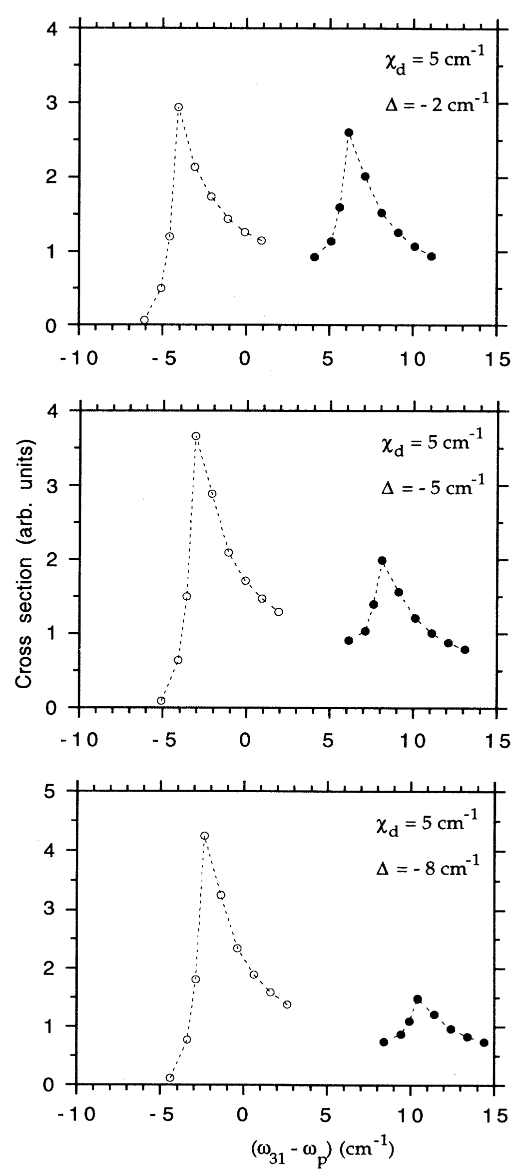


FIG. 4. Dependence of the bichromatic LICET line shape on strong-field detuning $\Delta=\omega_d-\omega_{32}$.

speed $v$ equal to an average value of $5\times10^4$ cm/s [1]. Data are given in arbitrary but comparable units.

In Fig. 3 the excitation spectrum calculated for $\Delta=-4$ $cm^{-1}$ and $\chi_d=2$, 4, and 6 $cm^{-1}$ is reported, showing the dependence of the line shape on the strong-field intensity. As expected from the preceding section, the spectrum exhibits two peaks separated by the quantity $\Omega=\sqrt{\Delta^2+4\chi_d^2}$. For a Rabi frequency $\chi_d=2$ $cm^{-1}$ (corresponding to a laser intensity $I_d\approx2$ $MW/cm^2$), the peak centered at $\omega_{31}-\omega_p=(\Omega+|\Delta|)/2$ (black dots), superimposed to the quasistatic wing of the peak centered at $\omega_{31}-\omega_p=(|\Delta|-\Omega)/2$ (white dots), is just visible. At increasing intensities the two features move far away, with one peak falling down in intensity and the other standing out.

In Fig. 4 the excitation spectrum calculated for $\Delta=-2$, $-5$, and $-8$ $cm^{-1}$ and $\chi_d=5$ $cm^{-1}$ (corresponding to a laser intensity $I_d\approx13$ $MW/cm^2$) is reported, showing the dependence of the line shape on the strong-field detuning $\Delta$. For $\Delta=-2$ $cm^{-1}$ the two peaks have almost the same intensity. At larger values of $\Delta$, the peak centered at $\omega_{31}-\omega_p=(|\Delta|-\Omega)/2$ (white dots) increases in intensity and moves toward the red side of the spectrum, as the peak centered at $\omega_{31}-\omega_p=(\Omega+|\Delta|)/2$ (black dots) moves toward the blue side of the spectrum, falling down in intensity. As far as $\Delta/\chi_d\gg1$, the effect of the dressing field disappears and the bichromatic LICET spectrum reduces to the weak-field monochromatic LICET spectrum.

The validity of the adiabatic approximation (3.9) has been checked by comparison with results obtained by direct numerical integration of Eq. (2.21). As expected from the discussion presented in Sec. III, the adiabatic approximation gives excellent results for collisions with large impact parameter $b$, while it becomes less accurate as far as $\Omega$ becomes comparable with the Fourier width of the collisional interaction $\Delta\omega\approx v/2\pi b$. However, since collisions with large $b$'s are responsible for the behavior of the line shape near the two peaks, it turns out that, in view of an accurate evaluation of the line shape, the adiabatic approximation is not required to be strictly valid for any possible collision, but only for collisions providing the main contribution to the spectrum in the detuning interval of interest.

In conclusion, we have shown that LICET studies performed in the bichromatic configuration can lead to a better understanding of the effect of a Stark dynamic shift of the atomic levels involved in the process on the line shape of the excitation spectrum. Calculations performed for the Eu-Sr system show that a detectable splitting of the peak (of the order of many $cm^{-1}$) is expected even for moderate intensities of the dressing field (of the order of a few $MW/cm^2$), making this configuration very promising for a quantitative experimental study of the collisional-radiative interaction dynamics typical of LICET processes in the strong-field regime. In light of the very recent results reported in Ref. [14], where strong-field monochromatic LICET calculations have been extended to include the magnetic degeneracy of states involved in the collision, we expect that the effect of sublevel degeneracy, ignored in the present model, might be a slight broadening of the line shapes reported in Figs. 3 and 4, leaving unaffected the main conclusions of the paper.

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