Inelastic fast-electron–hydrogen-atom collision in a laser field

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The laser-assisted electron-impact excitation of $2s$ and $2p$ dressed states of hydrogen is studied. The interaction of a linearly polarized laser field with the colliding system is treated by the nonperturbative Floquet theory, while the interaction of the fast incident electron with the target atom is treated in the first Born approximation. The nonresonant collision in a low-intensity, low-frequency laser field is dominated by the process with no exchanged photon with the field, while at larger intensities the collision accompanied by the exchange of real or virtual photons is much more important. For the stimulated bremsstrahlung, the $1s \rightarrow 2s$ transition is strongly influenced by the resonant coupling between the final $2s$ and the intermediate $3p$ states. In the resonant cases, where laser frequency matches a transition frequency between the initial or final and an intermediate state, the collisions with the largest cross section are those in which the projectile induces a dipole transition in the atom.

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I. INTRODUCTION

Interest in the problem of electron scattering by atoms in a laser field arises from its application to plasma heating by an intense electromagnetic field. Also, the use of a laser makes it possible to measure certain electron-atom scattering parameters that would not otherwise be accessible to experiment [1]. The conceptual interest in the laser modification of electron-atom collisions comes from a number of effects that are unobservable in the absence of the field.

Since the first observation of multiphoton processes in the laser-assisted collision [2], a considerable effort has been made to study these processes both experimentally and theoretically. In the experiments performed to date [3], the photon energy (CO_2 laser) was small compared with the electron energy. The experimental data concerning the large-angle scattering are in reasonable agreement with the Kroll-Watson-type approximations (KWA) [4], which neglect the internal degrees of freedom of the atom. However, recent experimental measurement [5] of the low-energy, small-angle elastic e-He collision clearly demonstrates that the KWA is inadequate for such conditions. The semiperturbative method [6,7] is valid for fast incident electrons and electric fields that are weak compared to the Coulomb binding field experienced by the atomic electrons, but can nevertheless be strong by laboratory standards. In this method, the interaction between the fast projectile and the target atom is treated perturbatively, by using the Born series. On the other hand, the laser-atom interaction is treated within first-order time-dependent perturbation theory, while the laser-projectile interaction is treated exactly, by using the Volkov wave function [8]. The semiperturbative theory exhibits a spurious divergence when the photon energy is close to the energy of an atomic transition. In such a resonant case, and/or at higher laser intensities when two or more atomic levels are brought into a resonance by the field, the nonperturbative techniques should be applied to treat the laser-atom interaction [9]. The non-Hermitian Floquet method [10–13] is appropriate if in a realistic experiment the intensity of the field is not too high, to prevent the target atom from decaying appreciably during the scattering experiment. Besides, the intensity must vary slowly over one optical cycle in order to describe the atom by a single Floquet wave vector varying adiabatically with the field strength before the collision occurred. This condition is usually fulfilled in a typical experimental situation.

In this paper, we study inelastic scattering, accompanied by the transfer of N photons, of fast incident electrons ($E_i = 500$ eV) by hydrogen atoms in the presence of a linearly polarized laser field. A detailed analysis is made of the transition between the dressed states $1s \rightarrow 2s$ and $1s \rightarrow 2p$ in hydrogen, in both the nonresonant and resonant cases [14]. Considering collisions at intermediate energies involving the excited states of the atom, one should note that the dynamics of these processes differs substantially from that involving the ground-state atom, due to closer coupling with other states, including those of the continuum spectrum [15].

We apply the nonperturbative first Born-Floquet (BF) theory [16], in which the laser field-(projectile-plus-atom) interaction is treated nonperturbatively, within the framework of the Floquet theory, while the interaction of the fast incident electron with the target atom is treated in the first Born approximation. The calculation is performed by expanding the wave functions of the target atom dressed by the field on a discrete basis of complex Sturmian functions, which allows us to take into account exactly the bound-continuum-state contributions, which is of crucial importance for electron impact excitation at intermediate energies. We neglect the exchange effects in the present work, since the field-free exchange effects are essentially negligible at the high impact energies considered here, and they are either smaller or slightly enhanced in the presence of a laser field [17].

II. THEORY

We consider the classical monochromatic and monomode linearly polarized field that is spatially homogeneous over atomic dimensions and has the electric field vector $\mathcal{E}(t)=\mathcal{E}_o \sin(\omega t+\phi)$ and the vector potential $\mathbf{A}(t)=\mathbf{A}_o \cos(\omega t+\phi)$, with $\mathbf{A}_o=c\mathcal{E}_o/\omega$.

The Hamiltonian of the electron-atom system in the presence of a laser field can be written

$$H=H_o+H_t+V_d \ , \tag{1a}$$

where

$$V_d=-\frac{e^2}{r_0}+\frac{e^2}{|\mathbf{r}_o-\mathbf{r}_1|} \tag{1b}$$

is the electron-atom interaction in the initial channel, and H_o and H_t are, respectively, the Hamiltonians of the unbound electron and of the atomic target in the presence of a laser field. The indices 0 and 1 refer to the projectile and the target electron coordinates, respectively. We adopt the velocity gauge to carry out the calculation [16].

The time-dependent Schrödinger equation for the nonrelativistic incident electron in a laser field can be exactly solved, giving the well known Volkov wave function

$$\chi_{\mathbf{k}}(\mathbf{r}_o,t)=(2\pi)^{-3/2}e^{i(\mathbf{k}\cdot\mathbf{r}_o-\mathbf{k}\cdot\boldsymbol{\alpha}(t)-E_k t/\hbar)} \ , \tag{2}$$

where $\boldsymbol{\alpha}(t)=\boldsymbol{\alpha}_o \sin(\omega t+\phi)$ and $\boldsymbol{\alpha}_o=e\mathcal{E}_o/m\omega^2$, and E_k is the energy of the unbound electron.

Further on, the time-dependent Schrödinger equation for the hydrogen atom in a laser field can be transformed, under the assumption that the intensity and the frequency are constant or vary adiabatically, by using the Floquet ansatz

$$\Phi(\mathbf{r}_1,t)=e^{-i\epsilon t/\hbar}\sum_{M=-\infty}^{\infty} e^{-iM\omega t}\mathcal{F}_M(\mathbf{r}_1) \ , \tag{3}$$

to the time-independent system of coupled equations for the harmonic components,

$$(\epsilon+M\hbar\omega-H_a)\mathcal{G}_M=V_+\mathcal{G}_{M-1}+V_-\mathcal{G}_{M+1} \ , \tag{4}$$

where H_a is the target Hamiltonian in the absence of the laser field, $V_\pm=\mp(e\hbar/2mc)\mathbf{A}_o\cdot\nabla_1$, and $\mathcal{F}_M=e^{-iM(\phi-\pi/2)}\mathcal{G}_M$. $\mathcal{G}_M$ is introduced in order to eliminate the phase ϕ from the Floquet equation (4) [16].

Since the atom is initially in the ground state and would ionize under the influence of the radiation field, the solutions of Eq. (4) are constrained by physical boundary conditions in coordinate space, namely, the harmonic components are regular at the origin, $r=0$, and behave as a superposition of outgoing waves at large distances, $r\to\infty$ [12,13]:

$$\mathcal{F}_M(\mathbf{r}_1)\xrightarrow[r_1\to\infty]{}\sum_{M'} f_{M'M}(\hat{\mathbf{r}}_1)r^{i\gamma_{M'}}\frac{e^{ik_{M'}r_1}}{r_1} \ , \tag{5}$$

where $\gamma_{M'}=me^2/\hbar^2k_{M'}$ and

$$k_{M'}=\left[\frac{2m}{\hbar^2}(\epsilon+M'\hbar\omega)\right]^{1/2} . \tag{6}$$

The branch of the square-root function in Eq. (6) is chosen so that the exponential function decreases at large distances in the closed channels and increases in the open channels, with an outgoing-wave behavior. The homogeneous system of equations (4), together with the boundary conditions, Eqs. (5) and (6), form an eigenvalue problem for the complex quasienergies ϵ,

$$\epsilon=\epsilon_i+\Delta_i-i\Gamma_i/2 \ , \tag{7}$$

where Δ_i is the shift from the unperturbed energy ϵ_i and Γ_i is the induced width.

The linear eigenvalue system of coupled equations (4) is solved by expanding the harmonic components on a discrete basis of complex Sturmian functions. By choosing such Sturmian functions, which oscillate and decrease exponentially at large distances, one can implement implicitly the boundary conditions in solving the system of homogeneous equations (4) [13,16].

The S-matrix element corresponding to the excitation of the dressed state Φ_f from the initial dressed state Φ_i is given, in the first Born approximation, by the expression

$$S_{f,i}^{B1,F}=\frac{-i}{\hbar}\int_{-\infty}^{+\infty}dt\,\langle\chi_{\mathbf{k}_f}(\mathbf{r}_o,t)\Phi_f^{(-)}(\mathbf{r}_1,t)|V_d \times|\chi_{\mathbf{k}_i}(\mathbf{r}_o,t)\Phi_i^{(+)}(\mathbf{r}_1,t)\rangle \ . \tag{8}$$

Here, $\Phi^{(+)}(\mathbf{r}_1,t)$ is the Floquet eigenstate corresponding to the outgoing-wave behavior. When expanded on the basis, it can be written [16]

$$\Phi^{(+)}(\mathbf{r}_1,t)=e^{-i\epsilon t/\hbar}\sum_{M=-\infty}^{+\infty}e^{-iM\omega t}e^{-iM(\phi-\pi/2)} \times\sum_{nl}c_{nl}^{(M)}r_1^{-1}S_{nl}^{\kappa}(r_1) \times Y_{lm_i}(\hat{\mathbf{r}}_1) \ , \tag{9}$$

where $c_{nl}^{(M)}$ are the coefficients of the expansion of the radial Floquet harmonic component on Sturmian basis. The $\Phi_f^{(-)}$ is the Floquet state which asymptotic behavior is to decrease exponentially in the closed channels and to increase exponentially in the open channels with ingoing-wave behavior. It can be obtained within the time-reversal operator formalism [16,18].

Since the energy widths of the initial and final states of the target are different, the time integration in Eq. (8) does not lead to an energy-conservation delta function. In order to make possible the numerical calculation, we neglect the width of the initial and final states in doing time integration in Eq. (8). This approximation is justified for the electron impact excitation of $n=2$ level of H in the field of Nd:YAG (where YAG denotes yttrium aluminum garnet) or higher-frequency lasers of low to moderate intensity treated in this work. The width of the initial state (usually very small) and that of the final state are not larger than 10^{-5} a.u. (hence the corresponding lifetime is of the order 10^{-12} s) for the intensity 10^{11} W/cm^2 and somewhat higher. Thus, the uncertainty of the photon energy related to the relevant virtual transitions does not exceed 0.1%, and the energy conservation would be violated very slightly by the above assumption.

Also, since the lifetime of the "dressed" hydrogen atom is much larger than the collision time (being of the order 10^{-16} s), the target would not disintegrate appreciably during the scattering experiment, and therefore the cross section is not affected significantly by neglecting the width of the initial and final states in Eq. (8).

From now on, working in atomic units (a.u.), we obtain after the integration on the time variable in Eq. (8)

$$S_{f,i}^{B1,F}=\frac{i}{2\pi}\sum_{N=-\infty}^{+\infty}\delta(E_{k_f}+\epsilon_f-E_{k_i}-\epsilon_i-N\omega)f_{f,i;N}^{B1,F} . \tag{10}$$

Here, $f_{f,i;N}^{B1,F}$ is the first Born approximation to the scattering amplitude for the transition between the dressed state $i\rightarrow f$ in hydrogen, with the transfer of N photons. It is given by

$$f_{f,i;N}^{B1,F}=\frac{-2}{K^2}e^{-iN\phi}\sum_{M,M'=-\infty}^{+\infty} i^{M'-M}J_{N-M'+M}(\mathbf{K}\cdot\boldsymbol{\alpha}_o) \times\langle\mathcal{F}_M^f|e^{i\mathbf{K}\cdot\mathbf{r}}-1|\mathcal{F}_{M'}^i\rangle . \tag{11}$$

In this equation, $\mathbf{K}=\mathbf{k}_i-\mathbf{k}_f$ is the momentum transfer in the collision, and J_l is the Bessel function. The calculation of the scattering amplitude is performed by expanding the plane wave $e^{i\mathbf{K}\cdot\mathbf{r}}$ and the Floquet harmonic components $\mathcal{F}_M$ onto spherical harmonics, and the subsequent expansion of $\mathcal{F}_M$ onto complex radial Sturmian functions. The numerical evaluation of the radial integrals involved is presented in Taïeb *et al.* [19].

The first Born-Floquet differential cross section given by

$$\frac{d\sigma_{f,i;N}^{B1,F}}{d\Omega}(\mathbf{K},\boldsymbol{\alpha}_o)=\frac{k_f}{k_i}|f_{f,i;N}^{B1,F}|^2 \tag{12}$$

does not depend on the phase ϕ of the laser field, due to the inability of the collision time to be defined, as a result of the approximation of the projectile wave packet by a monoenergetic beam of infinite duration [16].

When comparing theoretical results to the experimental cross sections, one first has to obtain cross sections over a fine mesh of intensities and then convolute them with a realistic spatiotemporal distribution of intensities of the laser beam. The ponderomotive acceleration of the projectile when penetrating and leaving the interaction region should also be taken into account, unless the pulse is extremely short.

III. RESULTS AND DISCUSSION

The nonperturbative BF theory with the Sturmian basis expansion takes fully into account the target atom distortion induced by a laser field. Such a distorted atom acts on the projectile by a long-range dipole potential ($\sim 1/r^2$), which is proportional to the (field-dependent) nondiagonal polarizability for the inelastic transition under study [20]. The long-range potential affects mainly the distant collisions, which contribute to near forward scattering. For collisions at larger scattering angles the target dressing becomes less important, and under nonresonant conditions one can model the atom by a structureless center of force. The laser-atom interaction may be considered as nonresonant if, for a given frequency, the intensity does not exceed a certain limit. The condition on the intensity is more stringent if the laser frequency is comparable to any characteristic atom excitation frequency. Thus, for a collision involving a $n=2$ level of hydrogen, in the field of an Nd laser, the condition is satisfied provided that $\mathcal{E}_o \leq 10^6$ V/cm [7].

The results presented in this paper are obtained for a geometry in which the polarization vector of the field $\mathcal{E}$ is parallel to the direction of the momentum transfer $\mathbf{K}$, varying thus with the scattering angle θ, and with the number of photons transferred in the collision. The reason we adopt this geometry is that the angular part of the scattering amplitude may be simplified and, what is more important, because for small momentum transfers, when approximately $\mathbf{k}_o\perp\mathbf{K}$, the coupling of the colliding system with the field has its minimal value for $\mathcal{E}_o\|\mathbf{k}_o$, and its maximal value for $\mathcal{E}_o\|\mathbf{K}$ [16,21]. Although in a realistic experiment the choice of the geometry $\mathcal{E}_o\|\mathbf{K}$ causes inconveniences because of the necessity of rotating the laser beam for each N and θ, the data concerning the experimental measurements of the elastic-electron–helium-atom collision for this geometry have been recently reported [5]. Moreover, in the case of small-frequency, small-momentum transfer collisions, the results referring to the geometry $\mathcal{E}_o\|\mathbf{K}$ should be very close to those obtained for $\mathcal{E}_o\perp\mathbf{k}_o$.

A. $1s\rightarrow 2s$ transition in the low-frequency laser field

We present in Fig. 1 the differential cross section for the electron-impact excitation of the $2s$ state of atomic hydrogen, as a function of scattering angle θ, for two fixed intensities $I_1=1.327\times10^9$ W/cm^2 (corresponding to $\mathcal{E}_o=10^6$ V/cm) and $I_2=7\times10^{11}$ W/cm^2 and for two "soft" laser photon energies $\omega_1=1.165$ eV (corresponding to $\lambda=1064$ nm) and $\omega_2=2$ eV. The results obtained with no target dressing (KWA) are also presented, for the intensity I_2. In light of the above discussion, we see that for the frequency ω_1, the intensity I_1 is in the domain of validity of a perturbation theory, in contrast to the intensity I_2, which is close to the four-photon ionization threshold of the $2s$ state [12]. We have not calculated the cross section above the threshold, since at higher intensities the $2s$ eigenenergy curve exhibits many avoided crossings with high Rydberg levels and the atom in this state decays too rapidly for the laser-assisted electron scattering experiment to be performed. We show in Fig. 2 the BF differential cross sections as a function of field intensity of a laser with the frequency ω_1, at scattering angle $\theta=0.5°$, and $\theta=10°$.

The general feature of a collision in the low-frequency, low-intensity laser field is the dominance of cross sections with no exchanged photon, $N=0$, and their indistinguishability for both low frequencies. The cross section is almost identical to the field-free cross section, since at low intensities the only nonnegligible contribution to the scattering amplitude, for $N=0$, gives the term that contains in both the initial and final states the Floquet harmonic components with photon index 0, $\mathcal{F}_0^{1s}$, and $\mathcal{F}_0^{2s}$,

which reduce to the nonperturbed atomic wave functions of $1s$ and $2s$ states, respectively, in the zero-field limit. For the same reason, the $N=0$ cross section is a least one order of magnitude larger than that corresponding to the exchange of one or two photons with the field. As the intensity increases, the $N=0$ cross section decreases slowly in the forward direction, while it decreases rapidly at a larger scattering angle considered here. This behavior is caused by (1) a smaller contribution of the zero-order Floquet component to the entire (normalized) Floquet wave function at higher than at lower intensity; (2) decrease of the component $M'=M=0$ of the scattering amplitude, Eq. (11), due to the decrease of Bessel function $J_0(\mathbf{K}\cdot\boldsymbol{\alpha}_0)$ as its argument increases in its first oscillation; (3) destructive interference of the virtual processes of no photon absorption and those in which the atom, after the interaction with the projectile, absorbs or emits a photon before ending in the final state. Near the zeros of $J_0(\mathbf{K}\cdot\boldsymbol{\alpha}_0)$, the above-mentioned virtual processes are of the same order of magnitude, and their interference causes the shift of the minima of the BF results with respect to those obtained in the KWA. As seen in Fig. 1, the target dressing has a relatively large effect on the $N=0$ cross section up to $\theta=25°$, while at larger angles the disagreement between BF and KWA differential cross sections persists in the region of minima. The influence of the target dressing diminishes at lower intensity.

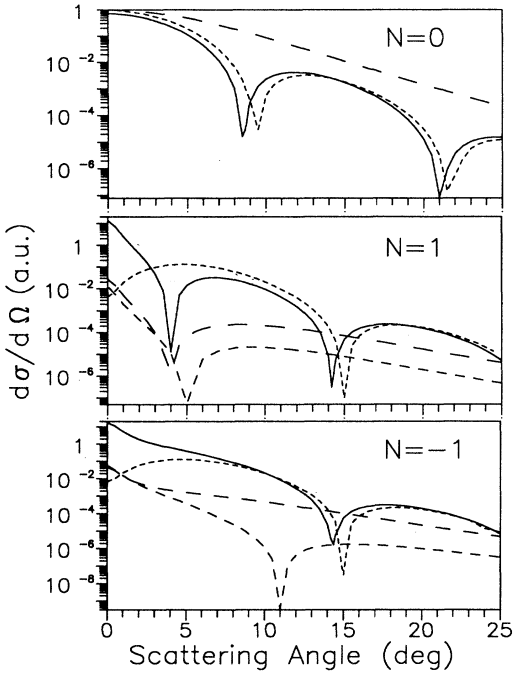

FIG. 1. Differential cross section (in atomic units) for the electron-impact excitation of the $2s$ state of atomic hydrogen in the presence of a linearly polarized laser field (parallel to the momentum transfer) as a function of the scattering angle θ (in deg). The incident electron energy is $E_i=500$ eV. N is the net number of photons exchanged by the e-H system and the field in the collision. Solid line: cross section corresponding to the laser intensity $I=7\times10^{11}$ W/cm^2 and $\omega=1.165$ eV. Dotted line: results obtained for the same field parameters, but with no target dressing. Long dashed line: $I=1.327\times10^9$ W/cm^2 and $\omega=1.165$ eV. Short-dashed line: the same but for $\omega=2$ eV.

At low intensity, $N=\pm1$ cross sections are two orders of magnitude larger for scattering in the forward direction than at $\theta=10°$, due to the strong S-P coupling at small scattering angles. Indeed, in the forward direction the dominant contribution to the $N=1(-1)$ cross section gives the terms corresponding to the absorption (emission) of one photon and a successive interaction with the projectile [i.e., $M'=1(-1)$, $M=0$], and vice versa [$M'=0$, $M=1(-1)$]. While the last term is slightly smaller than the former one for the $N=1$ cross section, it is 2–3 times larger in the case of the $N=-1$ cross section due to the strong resonant coupling of $2s$ and $3p$ states. The two terms interfere constructively in both cases. The situation is different at $\theta=10°$. The dominant contribution to the $N=1(-1)$ cross sections gives the term with $M=M'=0$ Floquet components, which is almost identical for both processes, and the several times smaller contribution of the term with $M'=0$, $M=1(-1)$. The mutual interference of these terms (constructive for $N=-1$, and destructive for $N=1$ cross section) may provide a relatively large effect on the cross section. Furthermore, the minima in the cross section $N=1(-1)$, at $\theta=4°$ for the frequency ω_1 and $\theta=5°$ (11°) for the frequency ω_2, are the consequence of an almost exact cancellation of terms corresponding to virtual processes with no photon absorption, and those in which the atom absorbs (emits) a photon in the initial state before interacting with the projectile, and vice versa [7].

As the intensity increases, cross sections for the transfer of $|N|\geq1$ photons increase rapidly at a fixed scattering angle, due to the increasing contribution of the first- and higher-order processes in the dressing of the target atom and the projectile by the laser field. Moreover, cross sections for the emission of photons are by a constant factor larger than the analogous ones for the absorption of photons, owing to the resonant coupling of the final and higher states. Indeed, if the intensity is not too high, the electron-impact excitation of hydrogen to a state of the $n=2$ level may be affected by resonant transition between the final and an intermediate states of higher energy in the case of stimulated bremsstrahlung, but not in the case of inverse bremsstrahlung. The physical explanation of this feature is that in the low-intensity limit the fast-electron–atom collision is governed by the second-order processes in which the atom interacts only once with both the projectile and the laser field (for more detailed analysis, see Sec. III C of this article and Ref. [7]). In addition, the initial $1s$ state could not be resonantly coupled with any other state by the absorption of a few low-energy photons. Therefore, a rough estimate of how much resonant coupling enhances the cross section for the laser-assisted $1s\rightarrow2s$ transition with the transfer of $N<0$ real photons could be made by considering the ratio of the cross section for the emission to the absorption of N photons. This ratio increases with N, becoming larger than ten for $|N|=3,4$ cross sections (not shown in

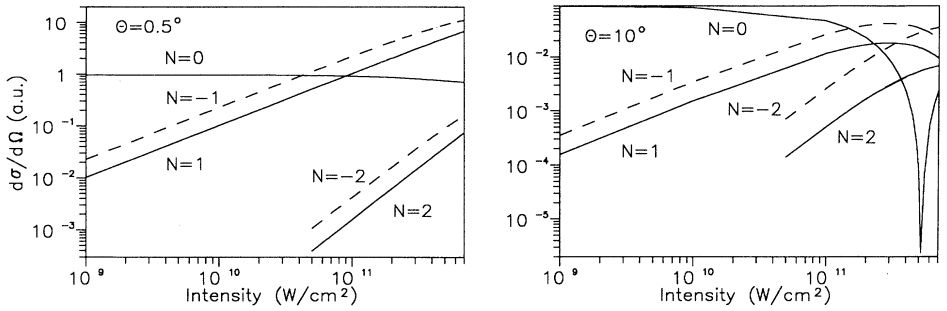

FIG. 2. Differential cross section for the electron-impact excitation of the dressed $2s$ state in the field of a varying intensity and frequency $\omega=1.165$ eV, at fixed scattering angles $\theta=0.5°$ and 10°. Solid line: results with Floquet dressing corresponding to $N=0,1,2$ absorbed photons. Dashed line: results for $N=-1,-2$ emitted photons.

Fig. 2), below the four-photon ionization threshold of the $2s$ state. This rise is pronounced by the three-photon resonant coupling of $2s$ state by a very high Rydberg level [12]. However, for all intensities below the small region near the threshold, the first-order resonant coupling between $2s$ and $3p$ states governs the $1s\rightarrow 2s$ transition in the case of stimulated bremsstrahlung. Furthermore, for a fixed scattering angle and a fixed intensity, the cross section for $|N|>1$ decreases rapidly with N. For such processes, the harmonic component of the target wave function that contributes the most to the cross section is that having photon number 0 or 1, depending on whether the scattering occurs at a large or small angle. The rest of the real photons are transferred to the projectile by the field. Intuitively, one can explain these features: (1) the Coulomb binding field shields the laser field on low atomic orbitals; (2) the projectile scattered by a neutral atom moves in faster decreasing potentials than the Coulomb one, so that the number of real photons transferred by the field in the collision could not be large. Since the effect of strong Coulomb potential on the scattering process would increase with decreasing impact parameters, one may expect that the multiphoton character of the laser-assisted collision should be stronger at large scattering angles.

We can see from Fig. 1 that the KWA is fairly inadequate for a description of forward scattering with the transfer of the $|N|=1$ photon, since this approximation does not take into account the target dressing. We have also compared the results of our nonperturbative first Born-Floquet method with those of semiperturbative approach, in which the target dressing is treated in first-order perturbation theory [6,7]. The semiperturbative amplitude could be calculated to arbitrary accuracy, directly from Eq. (11), by neglecting the terms with $|M|+|M'|>1$, and by replacing the involved Floquet harmonic components with their perturbative counterparts [16]. The results of both methods are almost identical up to the intensity 10^{10} W/cm^2. At higher intensities, the discrepancy of $N=1$ cross sections is within 10%. The largest difference between the results of the two methods is for the $N=-1$ cross section in the forward direction, and for the $N=0$ cross section at $\theta=10°$, in the region of the minimum at $I=5.2\times 10^{11}$ W/cm^2, up to the four-photon ionization threshold of $2s$ state. In this region, the behavior of these cross sections is determined by the complicated mechanism of interferences of components of the scattering amplitude. It is obvious that the first-order theory could not produce all the relevant components of the scattering amplitude with sufficient accuracy.

B. $1s\rightarrow 2p$ transition

We show in Fig. 3 the differential cross section for the $1s\rightarrow 2p$ transition, in a field of varying intensity and fre-

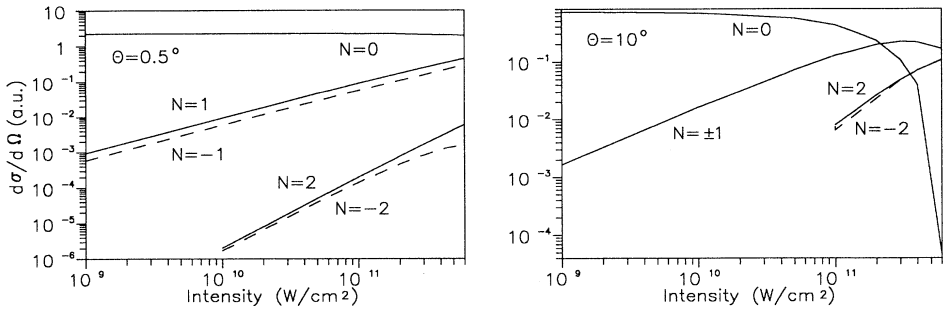

FIG. 3. Same as Fig. 2, but for the excitation of the $2p$ state of atomic hydrogen.

quency $\omega_1 = 1.165$ eV, at a fixed scattering angle $\theta = 0.5°$, or $\theta = 10°$, respectively. The limiting intensity shown in Fig. 3, $I = 6 \times 10^{11}$ W/cm^2, is near the four-photon ionization threshold of the $2p$ state.

The qualitative behavior of the cross section for the $1s \rightarrow 2p$ transition is, in general, similar to the analogous one corresponding to the $1s \rightarrow 2s$ transition. The most remarkable differences between the two transitions are the following.

(1) In the forward direction, the $N = 0$ cross section for the $1s \rightarrow 2p$ transition is larger than that for the $1s \rightarrow 2s$ transition, as in the field-free case.

(2) In the forward direction, the $N = \pm 1$ cross sections for the $1s \rightarrow 2p$ transition are at least one order of magnitude smaller than the analogous ones for the $1s \rightarrow 2s$ transition, the difference being smaller at larger intensities. At a fixed scattering angle and fixed intensity, cross sections for the absorption of one or two photons are larger than the analogous ones for the emission, contrary to the $1s \rightarrow 2s$ transition. The difference between $N = \pm 1$ cross sections is much smaller than for the $1s \rightarrow 2s$ transition. These features can be explained by the fact that the resonant coupling between the final $2p$ and higher intermediate states does not affect the $1s \rightarrow 2p$ transition as strongly as the analogous coupling of the final $2s$ state does in the case of the $1s \rightarrow 2s$ transition. Indeed, the contribution of the virtual process of absorption (emission) of a single photon is larger in the initial than in the final state for the $1s \rightarrow 2p$ transition. Furthermore, while in the forward direction the dominant components of the scattering amplitude corresponding to $N = \pm 1$ exchanged photons with the field interfere constructively in the case of the $1s \rightarrow 2s$ transition, these terms interfere predominantly destructively for the $1s \rightarrow 2p$ transition, at the intensities considered here.

C. Resonant collisions

Since the Born-Floquet theory treats the interaction of the external field with the target atom nonperturbatively, it is particularly useful to study resonant collision [16,22]. Here we assume that the initial (or final) state is represented by the *diabatic* state, along which the initial character of the state is preserved. This imposes the condition that the variation of the frequency is to be fast enough for the atom to jump diabatically across the energy gap at avoided crossing of the resonant levels [13].

As the first example, we show in Fig. 4 the differential cross section for the electron-impact excitation of $2s$ and $2p$ dressed states of H, in the vicinity of one-photon $1s$-$2p$ resonance. The scattering angle is $\theta = 0.5°$ and the intensity of the field $I = 10^{10}$ W/cm^2. At this intensity, the real parts of the adiabatic quasienergy curves ϵ_{1s} and ϵ_{2p} exhibit an avoided crossing at a wavelength 121.67 nm. For the collision with no exchanged photon, $N = 0$, the cross section corresponding to the transition $1s \rightarrow 2p$ is larger than that for the transition $1s \rightarrow 2s$, similarly to the scattering without the field at a small scattering angle. Obviously, in this case the projectile-atom interaction governs the scattering. Moreover, a specific feature of laser-assisted collisions is that the $N = 0$ cross section decreases when one approaches the resonant frequency, due

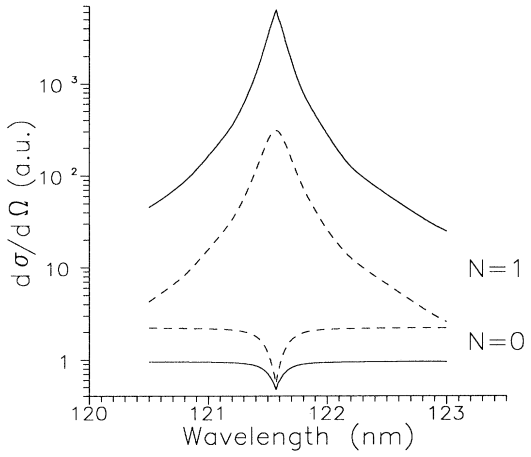

FIG. 4. Differential cross section in the vicinity of the one-photon $1s$-$2p$ resonance for the inelastic $1s \rightarrow 2s$ and $1s \rightarrow 2p$ electron hydrogen scattering in the presence of a linearly polarized laser field of intensity $I = 10^{10}$ W/cm^2, at a scattering angle $\theta = 0.5°$, vs wavelength (in nm). Solid line: cross section for the $1s \rightarrow 2s$ transition in H corresponding to $N = 0$ or 1. Dashed line: cross section for the $1s \rightarrow 2p$ transition in H, and $N = 0$ or 1.

to the smaller contribution of the $M = 0$ harmonic component to the entire Floquet wave function. Further, the $N = 1$ cross section for the $1s \rightarrow 2s$ transition is much larger than that for the $1s \rightarrow 2p$ transition, with a very pronounced maximum in the resonant region. This can be explained by the fact that the resonant collision is strongly dominated by the second-order process in which the atom is excited first to the $2p$ state by the absorption of a single photon, and then passes to the final dressed $2s$ state via the interaction with the projectile, which reduces to the dipole interaction at a small scattering angle. Indeed, the initial $1s$ state is prepared by the field on resonance approximately in a coherent superposition of the unperturbed $1s$ and $2p$ states. Thus, the $1s \rightarrow 2s$ transition accompanied by the absorption of one photon may be regarded as a field-free transition between "almost" unperturbed atomic states $2p$ and $2s$, which explains the sharp increase and the large maximum value of the $N = 1$ cross section in the resonant region. On the other hand, the $1s \rightarrow 2p$ transition is mainly affected by two second-order processes, in which the atom is excited to the final $2p$ state by the absorption of a single photon and a successive (or preceding) elastic scattering with the projectile in the final (or initial) state. The first term predominates over the second one, but it is still smaller than the term that gives the dominant contribution to the $1s \rightarrow 2s$ transition.

As the second example, we show in Fig. 5 the differential cross section for the $1s \rightarrow 2s$ and $1s \rightarrow 2p$ transitions in the vicinity of the one-photon $2s$-$3p$ resonance, at $\theta = 0.5°$ and $I = 1.327 \times 10^9$ W/cm^2. The couples of both $2s$-$3p$ and $2p$-$3s$ real adiabatic quasienergy curves exhibit avoided crossings close to the frequency 656.6 nm.

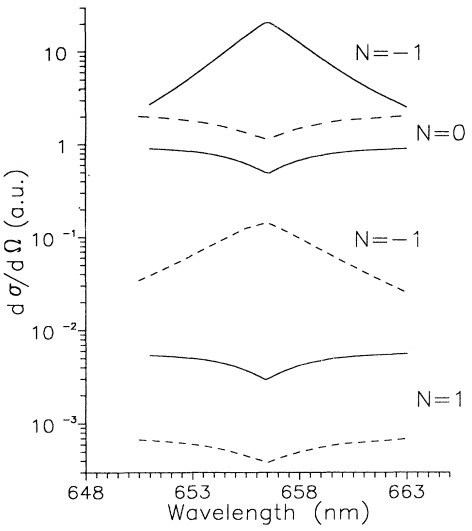

FIG. 5. Differential cross section in the vicinity of the one-photon $2s$-$3p$ resonance for the inelastic-electron–hydrogen scattering in the presence of a linearly polarized laser field of intensity $I=1.327\times10^{9}$ W/cm^2, at a scattering angle $\theta=0.5°$, vs wavelength (in nm). Solid line: cross section for the excitation of the $2s$ state and for $N=0,\pm1$. Dashed line: cross section for the excitation of the $2p$ state.

These resonances influence almost identically the $N=0$ cross section for both the $1s\rightarrow2s$ and $1s\rightarrow2p$ transitions, as the $1s$-$2p$ resonance does. The reason for such behavior is that the dominant contribution to the entire normalized Floquet wave function gives the harmonic component corresponding to the strong radiative transition between the resonant levels, so that the relative contribution of all other harmonic components vanishes close to the resonant frequency. Furthermore, in the forward direction the dominant contribution to the first Born scattering amplitude for the transfer of the $N=0,\pm1$ real photon is given by the second-order terms containing the Floquet harmonic component corresponding to the exchange of virtual photons whose number is equal to the number of real photons transferred in the collision. If the radiative process described by this Floquet component brings directly the initial (or final) state into resonance with an intermediate state, the collision is resonant with an increasing cross section in the resonant region; otherwise, it is nonresonant and the corresponding cross section decreases in this region. Our conclusion is confirmed by the behavior of the $N=-1$ cross sections for both inelastic transitions considered here, which increase when one approaches the resonant frequency, while the $N=0$ and $N=1$ cross sections decrease. We may also conclude from the above arguments that $1s\rightarrow2l$ low-frequency laser-assisted e-H collision is affected appreciably only by the resonant virtual dipole transitions ending in the final state, at low and moderate intensities [7]. Furthermore, the $N=-1$ cross section for the $1s\rightarrow2s$ transition is more than two orders of magnitude larger than that for the $1s\rightarrow2p$ transition. Both resonant collisions are governed by the process of excitation of the H atom by the projectile to a state with the principal quantum number $n=3$, and its successive deexcitation to a final state with $n=2$, by the stimulated emission of a photon. Since the small-angle collisions are dominated by S-P transitions, it is obvious that the $N=-1$ cross section for the $1s\rightarrow(3p)\rightarrow2s$ transition is much larger than that for $1s\rightarrow(3s)\rightarrow2p$.

We have also considered the $1s\rightarrow2s$ and $1s\rightarrow2p$ transitions in H in the vicinity of the two-photon $1s$-$2s$ resonance, at a fixed scattering angle $\theta=0.5°$ and the field intensity $I=1.327\times10^{9}$ W/cm^2. At this intensity, the real parts of the Floquet eigenenergies $1s$ and $2s$ undergo a true crossing at the resonant frequency. No structure is observed in the differential cross section for either of the laser-assisted inelastic e-H collisions studied in this work.

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[1] N. Andersen, J. W. Gallagher, and I. W. Hertel, Phys. Rep. **180**, 1 (1988).
[2] A. Weingartshofer, J. Holmes, G. Caudle, E. Clarke, and H. Kruger, Phys. Rev. Lett. **39**, 269 (1977).
[3] A. Weingartshofer, E. Clarke, J. Holmes, and C. Jung, Phys. Rev. Lett. **19**, 2371 (1979); A. Weingartshofer, J. Holmes, J. Sabbagh, and S. L. Chin, J. Phys. B **16**, 1805 (1983); B. Walbank, V. W. Connors, J. K. Holmes, and A. Weingartshofer, *ibid.* **20**, L833 (1987); B. Walbank, J. K. Holmes, and A. Weingartshofer, *ibid.* **20**, 6121 (1987); B. Walbank, J. K. Holmes, L. Le Blanc, and A. Weingartshofer, Z. Phys. D **10**, 467 (1988); B. Walbank and J. K. Holmes, Phys. Rev. A **48**, R2515 (1993).
[4] F. V. Bunkin and M. V. Fedorov, Zh. Eksp. Teor. Fiz. **49**, 1215 (1965) [Sov. Phys. JETP **22**, 844 (1966)]; N. M. Kroll and K. M. Watson, Phys. Rev. A **8**, 804 (1973); M. H. Mittleman, J. Phys. B **26**, 2709 (1993).
[5] B. Walbank and J. K. Holmes, J. Phys. B **27**, 1221 (1994).
[6] F. W. Byron, Jr. and C. J. Joachain, J. Phys. B **17**, L295 (1984); P. Francken and C. J. Joachain, Phys. Rev. A **35**, 1590 (1987); F. W. Byron, Jr., P. Francken, and C. J. Joachain, J. Phys. B **20**, 5487 (1987).
[7] P. Francken, Y. Attaourti, and C. J. Joachain, Phys. Rev. A **38**, 1785 (1988).
[8] D. V. Volkov, Z. Phys. **94**, 250 (1935).
[9] P. Francken and C. J. Joachain, Phys. Rev. A **41**, 3770 (1990); K. Unnikrishnan, Phys. Rev. A **38**, 2317 (1988).
[10] J. H. Shirley, Phys. Rev. **138**, B979 (1965).
[11] F. H. Faisal, *Theory of Multiphoton Processes* (Plenum, New York, 1987).

[12] R. M. Potvliege and R. Shakeshaft, Phys. Rev. A **40**, 3061 (1989).

[13] R. M. Potvliege and R. Shakeshaft, in *Atoms in Intense Laser Fields,* edited by M. Gavrila (Academic, New York, 1992).

[14] P. H. G. Smith and M. E. Flannery, J. Phys. B **24**, L489 (1991); **25**, 1021 (1992); S. Vučić, in *Proceedings of the 6th International Conference on Multiphoton Processes (ICOMP VI), Quebec City, Canada, 1993,* edited by D. K. Evans and S. L. Chin (World Scientific, New York, 1993).

[15] S. Vučić, R. M. Potvliege, and C. J. Joachain, J. Phys. B **20**, 3157 (1987); S. Vučić, R. M. Potvliege, and C. J. Joachain, Phys. Rev. A **35**, 1446 (1987).

[16] M. Dörr, C. J. Joachain, R. M. Potvliege, and S. Vučić, Phys. Rev. A **49**, 4852 (1994); Z. Phys. D **29**, 245 (1994).

[17] G. Ferrante, C. Leone, and F. Trombetta, J. Phys. B **15**, L475 (1982); F. Trombetta, C. J. Joachain, and G. Ferrante, *ibid.* **19**, 1081 (1986).

[18] M. Pont, R. M. Potvliege, R. Shakeshaft, and P. H. G. Smith, Phys. Rev. A **46**, 555 (1992).

[19] R. Taïeb, V. Véniard, A. Maquet, S. Vučić, and R. M. Potvliege, J. Phys. B **24**, 3229 (1991).

[20] H. R. J. Walters, Phys. Rep. **116**, 1 (1984).

[21] G. Ferrante, Phys. Rev. A **22**, 2529 (1980).

[22] Y. Gontier and M. Trahin, Phys. Rev. A **19**, 264 (1979); C. R. Holt, M. G. Raymer, and W. P. Reinhardt, Phys. Rev. A **27**, 2971 (1983).