Ionization of hydrogen by ion impact in the presence of a resonant laser field

A.B.Voitkiv[1,2], N.Toshima[2] and J.Ullrich[1]

[1] *Max-Planck-Institut für Kernphysik,*
Saupfercheckweg 1, D-69117 Heidelberg, Germany
[2] *Institute for Materials Science, University of Tsukuba,*
1-1-1 Tennodai, Tsukuba, Ibaraki 305 8573, Japan
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Abstract

We study the impact ionization of atomic hydrogen in collisions with fast ions assisted by a weak linearly polarized laser field whose frequency is resonant to the $1s$-$2p$ hydrogen transitions. We consider this process using a simple model in which the interaction between the atom and the resonant field is described in the Rotating-Wave approximation and the interaction of the field-dressed atom with the ion is treated within the Continuum-Distorted-Wave-Eikonal-Initial-State approach. Our consideration shows that the presence of the laser field can have a profound effect on all aspects of the impact ionization, including the angular and energy distributions of the emitted electrons, the total ionization cross section and the projectile scattering.

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I. INTRODUCTION

Ionization of atoms by charged projectiles like ions and electrons is one of the basic problems studied in atomic physics. The different aspects of electron emission in collisions with ions at relatively large impact energies, where the collision velocity v is higher than a characteristic orbiting velocity v_0 of the electron in its initial bound state, have been intensively studied, both experimentally and theoretically (for reviews see [1]-[6] where also the wealth of references to original works can be found).

There has been the growing interest to study ion-atom collisions assisted by laser fields [7]-[12]. In particular, the impact ionization of atoms and the charge transfer in fast and slow laser-assisted ion-atom collisions have been studied in some detail for the case when the frequency of the laser field was much smaller than typical frequencies for transitions from the ground atomic state. The inclusion of an external electromagnetic field into atomic collisions introduces new degrees of freedom and can, under certain conditions, very substantially influence the collision process.

In the present paper we address the question of how the process of atomic ionization by fast ionic projectiles can be modified in the presence of a laser field which is resonant with respect to atomic transitions. The field is assumed to be monochromatic, to have a linear polarization and be sufficiently weak so that the photoionization of the atom due to the laser field can be ignored.

The process of the laser-assisted impact ionization will be considered within a simple model which consists of two basic ingredients. Firstly, the interaction between the atom and the resonant laser field is described by using the Rotating-Wave approximation. Secondly, the interaction of the field-dressed atom with the ion is treated within the Continuum-Distorted-Wave-Eikonal-Initial-State approach.

Atomic units are used throughout except where otherwise stated.

II. GENERAL CONSIDERATION

Since the collision velocity is assumed to be relatively high we will use the impact parameter approximation. We shall neglect the (direct) action of the laser field on the atomic nucleus and the projectile. The nucleus of the atomic target is assumed to be at rest and is taken as the origin. The projectile moves along a straight-line trajectory $\mathbf{R}(t) = \mathbf{b} + \mathbf{v}t$. The electron coordinate with respect to the origin is $\mathbf{r}$. At the infinitely remote past ($t \to -\infty$) the atom is assumed to be in the ground state.

We shall suppose that the laser field is monochromatic and is switched on and off adiabatically at $t \to -\infty$ and $t \to +\infty$, respectively. We shall also assume that the dependence of the field on the space coordinates can be neglected, i.e. that the magnetic component of the field can be ignored and the electric component can be described using the dipole approximation,

$$\mathbf{F}(t) = \mathbf{F}_0 \cos(\omega_0 t + \varphi), \tag{1}$$

where ω_0 and φ are the field frequency and phase. The laser frequency ω_0 is supposed to be close to the $1s$-$2p$ transition frequency in hydrogen. The amplitude of the field strength, $|\mathbf{F}_0|$, is assumed to be much less than the typical atomic field $F_a = 1$ a.u..

A. Atomic states dressed by the laser field

We begin our consideration with building the initial and final atomic states in the presence of the laser field. These states are solutions of the Schrödinger equation

$$i\,\frac{\partial\psi_{i,f}(\mathbf{r},t)}{\partial t} = \left(\frac{\hat{p}^2}{2} - \frac{1}{r} + \mathbf{F}(t)\cdot\mathbf{r}\right)\psi_{i,f}(\mathbf{r},t), \tag{2}$$

where $\hat{\mathbf{p}}$ is the operator for the electron momentum and the interaction with the laser field is taken in the so called length gauge.

The initial state $\psi_i(\mathbf{r},t)$ satisfies the boundary condition $\psi_i(\mathbf{r},t\to-\infty)\to\chi_1(\mathbf{r})\exp(-i\varepsilon_1 t)$, where $\chi_1(\mathbf{r})$ is the ground state of the atom and $\varepsilon_1=-0.5$ a.u. is its energy. We shall look for this state by expanding it in the complete set of the free atomic states $\{\chi_n(\mathbf{r})\}$. Since the laser field is weak and its frequency is assumed to be quite close to the $1s$-$2p$ atomic transition frequency, the main contribution to the field-dressed state $\psi_i(\mathbf{r},t)$ in this expansion is given just by the ground, $\chi_1(\mathbf{r})$, and the first excited atomic states, $\chi_{2pm}(\mathbf{r})$, with $m=-1,0,1$. Therefore, one can write

$$\psi_i(\mathbf{r},t) = a_1(t)\chi_1(\mathbf{r})\exp(-i\varepsilon_1 t) + \sum_{m=-1}^{1} a_{2pm}(t)\chi_{2pm}(\mathbf{r})\exp(-i\varepsilon_2 t), \tag{3}$$

where $\varepsilon_2=-0.125$ a.u. is the energy of the excited $2p$ states. In (3) the sum runs over the magnetic quantum number m of the excited states and $g(t)$ and $a_{2pm}(t)$ are time-dependent coefficients to be determined.

For weak laser fields with $\omega_0\approx\varepsilon_2-\varepsilon_1$ the coefficients $g(t)$ and $a_{2pm}(t)$ can be found using the so called resonant (or rotating-wave) approximation (see e.g. [13]). Within this approximation the coefficients satisfy the following equations

$$\begin{aligned} i\frac{da_1(t)}{dt} &= \varepsilon_1 a_1(t) + \sum_{m=-1}^{1} W_{1,2pm}a_{2pm}(t)\exp(i\omega_0 t+\varphi_0) \\ i\frac{da_{2pm}(t)}{dt} &= \varepsilon_2 a_{2pm}(t) + W^{*}_{1,2pm}a_1(t)\exp(-i\omega_0 t-\varphi_0), \end{aligned} \tag{4}$$

where $W_{1,2pm}=0.5\langle\chi_1\mid\mathbf{F}_0\cdot\mathbf{r}\mid\chi_{2pm}\rangle$. The solution of (4) reads

$$\begin{aligned} a_1(t) &= \sqrt{\frac{\mid\Delta\mid+\Omega}{2\Omega}}\exp(-iEt) \\ a_{2pm}(t) &= \exp(-i\varphi_0)\frac{W_{2pm,1}}{\sqrt{0.5(\mid\Delta\mid+\Omega)\Omega}}\exp(-i(E+\omega_0)t). \end{aligned} \tag{5}$$

Here, $\Delta=\varepsilon_2-\varepsilon_1-\omega_0$ is the resonance detuning, $\Omega=\sqrt{\Delta^2+4|W_{1,2}|^2}$ is the Rabi frequency, $|W_{1,2}|^2=\sum_{m=-1}^{1}|W_{1,2pm}|^2$ and E is the real (quasi)energy of the field-dressed atomic state. The quasi-energy is given by $E=\varepsilon_1+0.5(\Delta-\Omega)$ for $\Delta>0$ and $E=\varepsilon_1+0.5(\Delta+\Omega)$ for $\Delta<0$.

In writing down the solutions (5) we have neglected the spontaneous radiative decay of the excited $2p$ states to the ground state and also the laser-induced coupling of these excited states to the atomic continuum. The main characteristics of these two processes are the widths of the excited states with respect to the spontaneous decay, Γ_r, and the photoionization, Γ_i.

The neglect of the laser-induced coupling to the continuum in the expressions (3) and (5) is only possible if $\mid W_{1,2}\mid\gg\Gamma_i$. Since $\mid W_{1,2}\mid\sim F_0$ and $\Gamma_i\sim F_0^2$ it is clear that the condition $\mid W_{1,2}\mid\gg\Gamma_i$ will always be fulfilled for not too intense laser fields. The necessary condition for the spontaneous decay to be ignored is $\mid W_{1,2}\mid\gg\Gamma_r$ which implies that the laser field should also be not too weak.

In order to justify the approximate solution (3) with the coefficients (5) the conditions $| W_{1,2} |\gg \Gamma_i$ and $| W_{1,2} |\gg \Gamma_r$ are necessary but not sufficient. Indeed, the spontaneous decay leads to the flow-out of the population from the excited states to the ground state. The laser-induced coupling to the continuum leads to the depopulation of the atomic bound states due to the photoeffect. In order to be able to neglect these effects the lifetimes of the excited states with respect to the photoeffect and radiative decay should be much larger than the duration of the laser pulse.

We shall return to the discussion of the above conditions when considering numerical examples in section III and for the moment simply assume that they are fulfilled.

In the final state the electron is in the continuum. If the atomic nucleus were absent the motion of the electron under the action of a laser field would be described by a Volkov solution [14]. In the non-relativistic limit this solution, given in the dipole approximation and the length gauge, reads

$$\begin{aligned}\psi_{\mathbf{k}}(\mathbf{r},t) &= \frac{\exp(i\mathbf{k}\cdot\mathbf{r})}{(2\pi)^3}\exp\left(i\frac{\mathbf{F}_0\cdot\mathbf{r}}{\omega_0}\sin(\omega_0 t+\varphi_0)\right)\exp\left(-\frac{i}{2}\int^t dt'\left(\mathbf{k}-\frac{\mathbf{F}_0}{\omega_0}\sin(\omega_0 t'+\varphi_0)\right)^2\right)\\ &= \frac{\exp\left(i\mathbf{k}\cdot\mathbf{r}-ik^2t/2-iU_pt\right)}{(2\pi)^3}\times\\ &\quad\sum_{p_1,p_2,p_3=-\infty}^{+\infty}\exp(-i\pi p_2/2)J_{p_1}(\alpha)\,J_{p_2}(\beta)\,J_{p_3}(\gamma)\,\exp(i(p_1+p_2+2p_3)(\omega_0 t+\varphi_0)),\end{aligned} \tag{6}$$

where J_p are the Bessel functions (see e.g. [15]), $\alpha = \mathbf{F}_0\cdot\mathbf{r}/\omega_0$, $\beta = \mathbf{F_0}\cdot\mathbf{k}/\omega_0^2$, $\gamma = F_0^2/(8\omega_0^3)$, $U_p = F_0^2/(4\omega_0^2)$ is the ponderomotive energy, $\mathbf{k}$ is the momentum of the electron and the sums run over integer numbers p_1, p_2, p_3. The Volkov solution differs from a simple plane wave by the presence of the additional time- and space-dependent factors in (6). The actual difference between the Volkov solution and the corresponding plane wave depends on the magnitudes of the quantities α, β and γ and, according to the properties of the Bessel functions, becomes quite small when these quantities are much smaller than 1.

For laser fields of interest for the present study the field frequency is relatively high ($\omega_0 \approx 3/8$ a.u.) but the field intensity is quite low ($F_0 \ll 1$) and we immediately observe that $U_p \ll 1$ and $\gamma \ll 1$. In addition, for not too large values of k one has $\beta \ll 1$. Besides, one should take into account that when considering bound-free electron transitions the values of the coordinate r are effectively restricted to $r \lesssim a_B$, where a_B is the typical dimension of the electron bound state. This means that $\gamma \ll 1$ as well.

In reality the target nucleus does influence the electron motion in the continuum and the Volkov state (6), where this influence is fully ignored, cannot be considered as a reliable approximation for the continuum states with not very high energies. However, the consideration of the previous two paragraphs can still be safely used to make a qualitative conclusion about importance (or unimportance) of the laser field for the continuum states. Therefore, the above observations, that the quantities α, β, γ and U_p are very small, together with that, what had been said about the field-induced resonance between the atomic continuum and the excited states, enables us to conclude that the effect of the laser field on the final electron state is very weak and can be ignored.

B. Description of the ion-atom interaction

The transition amplitude for the field-assisted impact ionization is given by

$$A_{\mathbf{k}}(\mathbf{b}) = -i\int_{-\infty}^{+\infty} dt\langle\Psi_f^{(-)}(\mathbf{r},t) \mid \left(\hat{H}_c - i\frac{\partial}{\partial t}\right) \mid \Psi_i^{(+)}(\mathbf{r},t)\rangle. \tag{7}$$

Here, $\mathbf{k}$ is the final momentum of the electron with respect to the atomic nucleus. Further,

$$\hat{H}_c = \frac{\hat{p}^2}{2} - \frac{1}{r} - \frac{Z_p}{s} \tag{8}$$

is the time-dependent Hamiltonian describing the motion of the electron in the combined field of the target nucleus and the projectile. The 'in' and 'out' states, $\Psi_i^{(+)}(\mathbf{r},t)$ and $\Psi_f^{(-)}(\mathbf{r},t)$, will be approximated using the Continuum-Distorted-Wave-Eikonal-Initial-State (CDW-EIS) approach [16] which has been proved to represent a rather simple but very valuable tool to treat ionization in fast ion-atom collisions. According to the CDW-EIS approximation the 'in' and 'out' states are taken as follows

$$\begin{aligned} \Psi_i^{(+)}(\mathbf{r},t) &= \psi_i(\mathbf{r},t) \exp\left(i\nu_p \ln(vs + \mathbf{v}\cdot\mathbf{s})\right) \\ \Psi_f^{(-)}(\mathbf{r},t) &= \chi_{\mathbf{k}}^{(-)}(\mathbf{r}) \exp(-ik^2/2t) \exp(\pi\nu_p/2)\, \Gamma(1+i\nu_p)\ {}_1F_1\left(-i\nu_p, 1, -ivs + i\mathbf{v}\cdot\mathbf{s}\right) . \end{aligned} \tag{9}$$

In (9) $\psi_i(\mathbf{r},t)$ is determined by Eqs.(3) and (5) and $\chi_{\mathbf{k}}^{(-)}(\mathbf{r})$ is the two-body continuum state of the free atom. This state satisfies the appropriate Coulomb boundary conditions and describes an electron which moves with the asymptotic momentum $\mathbf{k}$ with respect to the atomic nucleus ($k = \mid \mathbf{k} \mid$). Further, $\nu_p = Z_p/v$, Γ is the gamma-function and ${}_1F_1(a,b,z)$ is the degenerate hypergeometric function (see e.g. [15]).

When calculating cross sections it is more convenient to work with the transition amplitude in the momentum space, $S_{fi}(\mathbf{Q})$. The latter is related to the amplitude (7) by the two-dimensional Fourier transformation

$$S_{\mathbf{k}}(\mathbf{Q}) = \frac{1}{2\pi} \int d^2\mathbf{b}\, A_{\mathbf{k}}(\mathbf{b}) \exp(i\mathbf{Q}\cdot\mathbf{b}) \tag{10}$$

and the two-dimensional vector $\mathbf{Q}$ has the meaning of the transverse part, $\mathbf{Q}\cdot\mathbf{v} = 0$, of the total momentum transfer $\mathbf{q}$ to the target. Note also that the following identity (the Parceval theorem) holds

$$\int d^2\mathbf{b} \mid A_{\mathbf{k}}(\mathbf{b}) \mid^2 \equiv \int d^2\mathbf{Q} \mid S_{\mathbf{k}}(\mathbf{Q}) \mid^2 . \tag{11}$$

C. Evaluation of the cross sections

Compared to field-free collisions, the evaluation of cross sections in field-assisted collisions have a few special points. These points will be discussed in this subsection considering, as an example, the cross section differential in energy and solid angle of the emitted electron. In the field-free collisions such a cross section is known to yield most detailed information about the electron emission spectra and one can expect that it can provide valuable information also in the case of field-assisted collisions.

The relation between this differential cross section and the transition amplitude is given by

$$\frac{d^3S}{d\varepsilon_k d\Omega_{\mathbf{k}}} = \int d^2\mathbf{Q} \mid S_{\mathbf{k}}(\mathbf{Q}) \mid^2, \tag{12}$$

where $\varepsilon_k = k^2/2$ is the electron emission energy and $d\Omega_{\mathbf{k}} = \sin\theta_{\mathbf{k}} d\theta_{\mathbf{k}} d\varphi_{\mathbf{k}}$ with $\theta_{\mathbf{k}}$ and $\varphi_{\mathbf{k}}$ being the polar and azimuthal emission angles, respectively. We shall suppose that the projectile moves along the z-axis and, thus, the angle $\theta_{\mathbf{k}} = 0$ points to the direction of $\mathbf{v}$.

In the presence of the laser field the state $\psi_i(t)$ is a coherent superposition of the ground and excited atomic states. Therefore, the total transition amplitudes (7) and (10) are given by the sum of the corresponding amplitudes for the transitions from the ground and excited states. Since these transitions can lead to the population of the same states of the electron in the continuum it might seem that these 'partial' transition amplitudes will add up coherently in the cross section and, as a result, that there will be an interference effect between the emission from the ground and excited states.

The above interference, however, could arise only if the laser and projectile beams are specially adjusted in time on a time scale of (substantially) less than ω_0^{-1}. Since this is hardly possible, the cross section (12) still has to be averaged over the initial phase of the laser field. Indeed, taking in the consideration the projectile trajectory as $\mathbf{R}(t) = \mathbf{b} + \mathbf{v}t$ for *all* projectile-target pairs implies that the time of the closest projectile-target distance is given by $t \approx 0$ for *all* collisions. In reality, different projectile-target collisions take place at different times. This fact, which obviously makes no difference for the treatment of field-free collisions, has to be taken into account in the field-assisted collisions because different projectile-target collisions occur at different phases of the electromagnetic field. If we make here a natural assumption that the time for a projectile to traverse a gas of the target atoms is much larger than the period of the field oscillation, $T = \frac{2\pi}{\omega_0}$, then, in order to get the meaningful cross section, the "cross section" (12) still has to be averaged over the phase of the electromagnetic field. This can be done by formally considering that the time of the closest approach of a projectile to a target is the same for all projectile-target pairs but that the initial phase φ_0 is different for different collisions. Considering that all φ_0 from the interval $0 \leq \phi \leq 2\pi$ can be encountered in the collisions with the same probability, the cross section is obtained by averaging according to

$$\frac{d^3\sigma}{d\varepsilon_k d\Omega_{\mathbf{k}}} = \frac{1}{2\pi}\int_0^{2\pi} d\varphi_0 \, \frac{d^3S(\varphi_0)}{d\varepsilon_k d\Omega_{\mathbf{k}}}. \tag{13}$$

After the averaging procedure the cross section reads

$$\frac{d^3\sigma}{d\varepsilon_k d\Omega_{\mathbf{k}}} = \frac{\mid \Delta \mid + \Omega}{2\Omega} \times \frac{d^3\sigma_1}{d\varepsilon_k d\Omega_{\mathbf{k}}} + \frac{2 \mid W_{1,2} \mid^2}{(\mid \Delta \mid + \Omega)\Omega} \times \frac{d^3\sigma_2}{d\varepsilon_k d\Omega_{\mathbf{k}}}. \tag{14}$$

The first term in (14) describes the emission from the ground state while the second term represents the emission from the coherent superposition of the excited atomic states. The relative contributions of the excited bound states with different magnetic quantum numbers to the second term depend on a choice of the quantization axis and the collision geometry, in particular, on how the vector $\mathbf{F}_0$ of the laser electric field is directed with respect to the projectile velocity $\mathbf{v}$. In the full cross section these two partial cross sections are weighted with the factors which depend on the frequency and intensity of the laser field and can be conveniently expressed by introducing the reduced detuning $\eta = \frac{|\Delta|}{2|W_{1,2}|}$

$$\begin{aligned} \frac{\mid \Delta \mid + \Omega}{2\Omega} &= \frac{\eta + \sqrt{1+\eta^2}}{2\sqrt{1+\eta^2}} \\ \frac{2 \mid W_{1,2} \mid^2}{(\mid \Delta \mid + \Omega)\Omega} &= \frac{0.5}{(\eta + \sqrt{1+\eta^2})\sqrt{1+\eta^2}}. \end{aligned} \tag{15}$$

Note that in no case the first/second factor can become smaller/larger than 0.5.

For a weak laser field the magnitude of the quasi-energy E is always quite close to the energy ε_1 of the undisturbed atomic ground state. Therefore, in calculations of the first partial cross section, $\frac{d^3\sigma_1}{d\varepsilon_k d\Omega_{\mathbf{k}}}$, one can safely set $E = \varepsilon_1$. The quasi-energy $E + \omega_0$ may in general noticeably differ from the energy ε_2 of the undisturbed excited states. However, when the latter happens, the second factor in (14) becomes very small making the contribution of the excited states to the cross section negligible. Therefore, in the calculation of $\frac{d^3\sigma_2}{d\varepsilon_k d\Omega_{\mathbf{k}}}$ one can make the replacement $E + \omega_0 \to \varepsilon_2$.

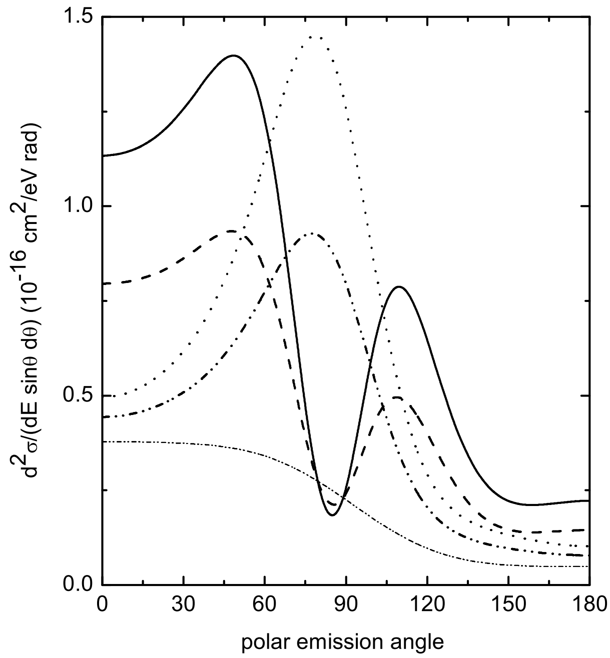

FIG. 1: Hydrogen ionization by 1 MeV/u C^{6+}. The doubly differential ionization cross section $d^2\sigma/(d\epsilon_k sin\theta_k d\theta_k)$ given as a function of the polar emission angle for a fixed emission energy of 0.1 a.u.. Solid and dash curves: results for $\eta = 0$ and $\eta = 0.5$, respectively, in the case of $\mathbf{F}_0 \parallel \mathbf{v}$. Dot and dash-dot curves: results for $\eta = 0$ and $\eta = 0.5$, respectively, in the case of $\mathbf{F}_0 \perp \mathbf{v}$. Dash-dot-dot curve: results for the field-free collision.

As obvious from the above consideration the structure of the other differential and total ionization cross sections in the presence of a resonant laser field is pretty similar to that of Eq.(14).

III. RESULTS AND DISCUSSION

In this section we consider the influence of a weak resonant laser field on the impact ionization of hydrogen by fast carbon nuclei C^{6+}. Throughout the section it is assumed that the strength F_0 of the laser field is fixed at 10^{-4} a.u. which corresponds to the laser field intensity of 3.5×10^8 W/cm^2.

At this intensity we obtain that $\mid W_{1,2} \mid \approx 4 \times 10^{-5}$ a.u. which is much larger than the widths of the $2p$ states with respect to the photoeffect $\Gamma_i \sim 10^{-9}$ a.u. and also than the spontaneous width of these states $\Gamma_r \sim 10^{-8}$ a.u.. Thus, the conditions $\mid W_{1,2} \mid \gg \Gamma_r$ and $\mid W_{1,2} \mid \gg \Gamma_r$, are very well fulfilled.

At such intensity the lifetimes of the atomic $2p$ states with respect to the photoionization are of about 10^{-8} s. The lifetime of these states with respect to the spontaneous radiative decay is about 10^{-9} s. Therefore, the effects of the photoionization and the spontaneous decay on the collision can be safely neglected provided the duration of the laser pulse is say of the order of 10^{-10} s or less.

After the laser pulse is switched on, the adiabatic dressed state (3) needs some time to be formed. The typical 'building time' τ for this state can be roughly estimated as the inverse of the

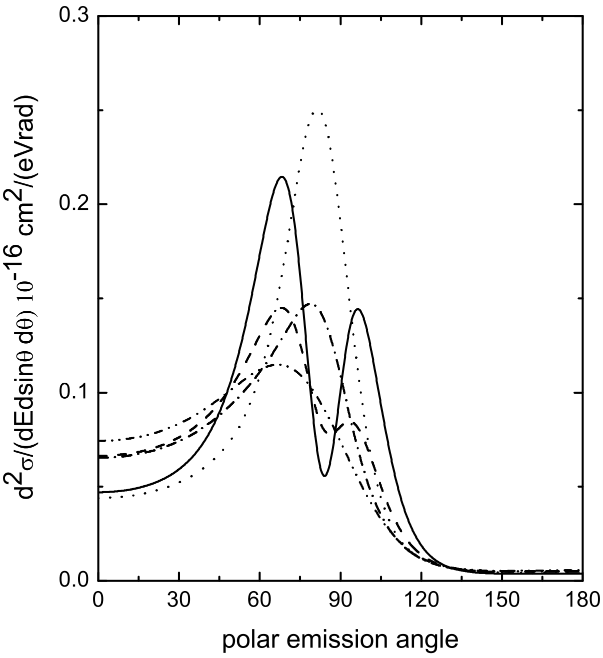

FIG. 2: Hydrogen ionization by 1 MeV/u C^{6+}. The doubly differential ionization cross section $d^2\sigma/(d\epsilon_k sin\theta_k d\theta_k)$ given as a function of the polar emission angle for a fixed emission energy of 0.5 a.u.. Solid and dash curves: results for $\eta = 0$ and $\eta = 1$, respectively, in the case of $\mathbf{F}_0 \parallel \mathbf{v}$. Dot and dash-dot curves: results for $\eta = 0$ and $\eta = 1$, respectively, in the case of $\mathbf{F}_0 \perp \mathbf{v}$. Dash-dot-dot curve: results for the field-free collision.

Rabi frequency, $\tau = 1/\Omega$, and in the resonance case $\tau = 1/\mid W_{1,2} \mid$. This time, obviously, has to be much smaller than the duration of the laser pulse. For the field with $F_0 = 10^{-4}$ a.u. one has $\tau \approx 6.5 \times 10^{-13}$ s and thus the pulse duration should not be shorter than, say, 5×10^{-12} s. Both this and the previous condition on the laser pulse duration are quite compatible with each other and can, for instance, be simultaneously fulfilled for the pulse duration of 10^{-10} s.

For the chosen parameters of the laser field one has $\alpha \sim 10^{-3}$-10^{-4}, $\beta \sim 10^{-3}k$, $\gamma \sim 10^{-8}$ and $U_p \sim 10^{-8}$ a.u.. Thus, α, γ and U_p are very small. Besides, for values of the electron momentum k, say, less than 10 a.u. β is very small as well. Taking into account that below the emission of electrons with high energies will not be considered we may conclude that the influence of the laser field on the electron states in the continuum can indeed be ignored.

As a final preliminary remark we note that the resonant laser field acting on a gas of the target atoms can produce a noticeable amount of almost mono-energetic photoelectrons having energies very close to $\varepsilon_2 + \omega_0$ which could be detected in experiment. According to the above estimates, however, the averaged space density of these electrons is much smaller compared to the density of the bound electrons in the target gas and, hence, they are not expected to influence the impact ionization. Since these electrons have neither the direct relation to nor the effect on the laser-assisted impact ionization, below they will not be considered.

We start our discussion of the laser-assisted impact ionization with considering the doubly differential ionization cross section, $\frac{d^2\sigma}{d\varepsilon_k \sin\theta_{\mathbf{k}} d\theta_{\mathbf{k}}}$, which is obtained by integrating the cross section (14) over the azimuthal angle of the emitted electron. Results for this cross section are shown in

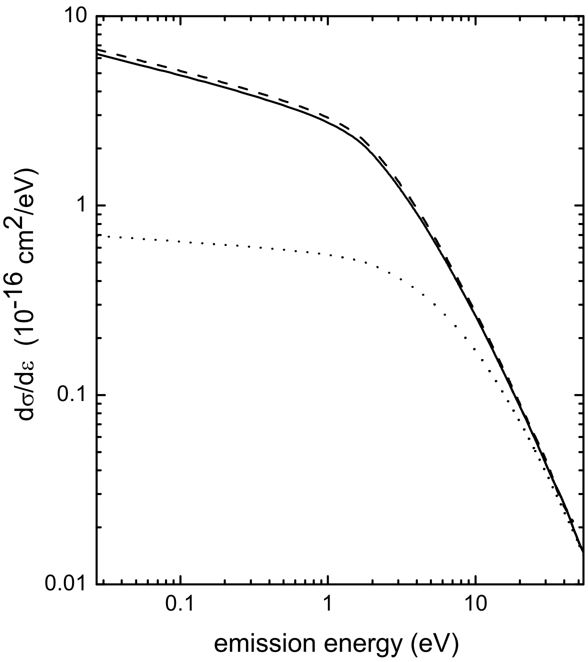

FIG. 3: Hydrogen ionization by 1 MeV/u C^{6+}. The energy spectrum of the emitted electrons, $d\sigma/d\varepsilon_k$. Solid and dash curves: results for $\mathbf{F}_0 \parallel \mathbf{v}$ and $\mathbf{F}_0 \perp \mathbf{v}$, respectively, and zero detuning. Dot curve: results for the field-free collisions.

figures 1 and 2. In these figures the cross section is given as a function of the polar emission angle for two electron emission energies (0.1 a.u. in figure 1 and 0.5 a.u. in figure 2) and different values of the reduced detuning $\eta = \frac{|\Delta|}{2|W_{1,2}|}$.

Two different collision geometries are considered in figures 1 and 2 with the field vector $\mathbf{F}_0$ either (i) directed along the collision velocity $\mathbf{v}$ or (ii) lying in the plane perpendicular to it. In our calculations the velocity axis was chosen as the quantization axis for the electron states.

It is seen in figures 1 and 2 that the field can have quite substantial effects on the doubly differential cross section. These effects arise due to the mixing of the ground and excited states by the resonant laser field and, naturally, reach their maximal magnitude when the resonance detuning is zero.

Different positions of the field vector $\mathbf{F}_0$ with respect to the collision velocity $\mathbf{v}$ lead to different populations of the excited atomic states. The emission pattern produced in collisions with atoms in excited states in general substantially depends on the magnetic quantum numbers of these states. As a result, besides the dependence on the resonance detuning, the shape of the angular emission spectrum is strongly affected by the collision geometry.

For instance, when the field vector $\mathbf{F}_0$ is directed along the collision velocity, the field couples the ground state only with the $2p_0$ excited state. The first Born approximation for the (laser field-free) ion-atom collisions predicts that, because of the selection rules, the electron emission under the impact ionization of the $2p_0$ state in fast collisions is suppressed in the directions perpendicular to the collision velocity. The higher order terms in the projectile-electron interaction, which are partly included in the CDW-EIS model, can substantially modify the emission pattern leading to

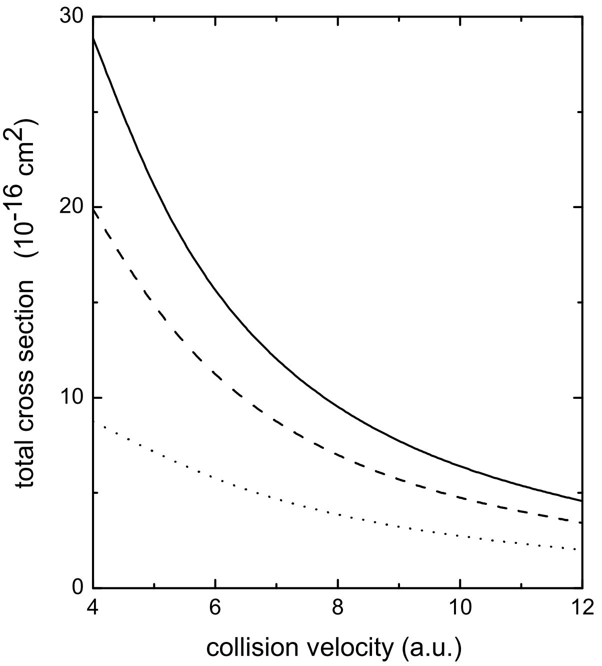

FIG. 4: Hydrogen ionization by C^{6+}. The total cross section given as a function of the collision velocity, $\mathbf{F}_0 \parallel \mathbf{v}$. Solid and dash curves: results for $\eta = 0$ and $\eta = 0.5$, respectively. Dot curve: results for the field-free collision.

the shift of the electron emission to the forward angles caused by the mutual attraction of the projectile and the electron in the final channel (the so called post-collision effect). However, the above suppression survives in the CDW-EIS model being somewhat altered by the post-collision effect.

In the field-assisted collisions, if the laser coupling of the $1s$ and $2p_0$ is sufficiently effective (as it is the case for $\mathbf{F}_0 \parallel \mathbf{v}$), there will be clear signatures of the emission from the $2p_0$ state which we observe in figures 1 and 2 as quite profound a minimum in the emission spectrum at angles close to 90^0.

For the other case of the collision geometry, considered in these figures, the laser field mixes the ground state only with the $2p_1$ and $2p_{-1}$ states. Due to the selection rules the emission produced in collisions with atoms in these states favors the directions perpendicular to the collision velocity. Therefore, instead of a minimum, for the second type of the collision geometry we observe a maximum in the emission spectrum at angles close to $\theta_k = 90^0$ (see figures 1 and 2).

In figure 3 we display the energy spectrum of the emitted electron. The effect of the resonant field is most prominent at the low emission energies. This reflects the fact that, compared to the impact ionization from the ground state, the population of the low-energy continuum is much more probable from the excited bound states. When the emission energy increases the difference between the energy spectra, produced in the field-free and field assisted collisions, decreases and becomes almost indistinguishable at sufficiently high emission energies. This can be understood by taking into account that the high-energy continuum states are populated mainly in collisions with relatively large momentum transfers when the atomic bound electron can be viewed as almost quasi-free independently of whether the electron is initially in the ground or in the excited states.

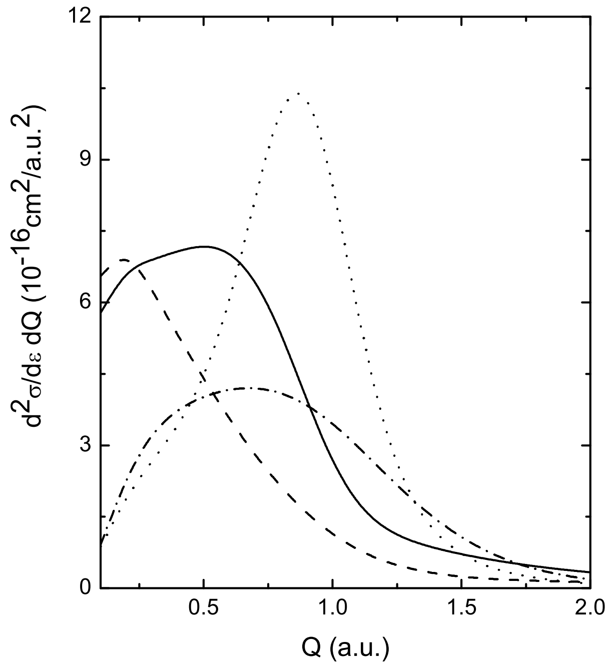

FIG. 5: Hydrogen ionization by 1 MeV/u C^{6+}. The doubly differential ionization cross section $d^2\sigma/(d\varepsilon_k dQ)$ given as a function of the transverse momentum transfer for a fixed emission energy of 10 eV, $\mathbf{F}_0 \parallel \mathbf{v}$ and $\eta = 0$. The solid and dash curves display results for the field-assisted and field-free collisions, respectively. Besides, dot and dash-dot curves show results for the field-assisted and field-free collisions, respectively, obtained by neglecting the interaction between the projectile and the target nucleus.

Therefore, the ionization arising from such collisions is on average not very sensitive to the presence of the laser field.

Some ideas about the overall effect of the laser field on the impact ionization can be obtained from figure 4 where the total ionization cross section is plotted as a function of the collision velocity. According to the figure the total ionization cross section can be strongly enhanced by the presence a weak resonant field. The reason for this enhancement is that, compared to the ionization from the ground state, the impact ionization from the excited states has larger total cross sections. When the collision velocity increases the cross sections for the impact ionization from the excited states decrease somewhat faster compared to that from the ground state. As a result, the field-assisted-to-field-free cross section ratio decreases with increasing the collision velocity.

The presence of the resonant laser field changes not only the electron emission spectra and the total ionization cross section but also alters the projectile scattering. This can be seen in figure 5, where the cross section $d^2\sigma/d\varepsilon_k dQ$, which is differential in the electron emission energy and the transverse part of the momentum transfer to the target, is displayed. In contrast to the electron emission spectra, integrated over the transverse momentum transfer, the cross section differential in Q may be substantially influenced by the interaction between the projectile and the target nucleus and, therefore, the latter has to be taken into account in the treatment. In our calculation this was done by employing the symmetric eikonal approximation in which the interaction between the ion and the atomic nucleus is included in the initial and final wavefunctions by introducing eikonal phases (similar to that entering the first equation in (9)).

It follows from the results, shown in figure 5, that both the laser field and the interaction between the projectile and the target nucleus may have substantial effects on the projectile scattering. Besides, it worth noting that the laser field and the interaction with the target nucleus counteract. The interaction with the target nucleus tends to increase the projectile scattering to smaller angles (which is a well known effect, see e.g. [17]). In contrast, the modification in the electron state in the atom, caused by the presence of the laser field, affects the interaction between the projectile and the electron in such a way which leads to the shift of the cross section to larger values of Q and, thus, the projectile scattering to larger angles.

IV. CONCLUSIONS

In this article we have attempted to address the question of the modification of the ion-impact ionization of atoms by the presence of a resonant laser field. As the simplest theoretical example of the field-assisted impact ionization we have considered in some detail the ionization of hydrogen in the presence of a linearly polarized laser field which is resonant to the $1s$-$2p$ hydrogen transitions.

We have treated this process by developing a simple model in which the interaction between the atom and the resonant field is described in the Rotating-Wave approximation and the interaction of the field-dressed atom with the ion is taken into account within the Continuum-Distorted-Wave-Eikonal-Initial-State approach.

The intensity of the laser field was chosen to be weak enough in order that the ionization of hydrogen due to photoeffect can be ignored. Besides, the weakness of the laser field allowed us to neglect its influence on the continuum states of the electron and also the presence of the non-resonant atomic bound states. However, due to the resonance condition, $\omega_0 \approx \varepsilon_2 - \varepsilon_1$, even quite a weak laser field can very effectively couple the ground and excited atomic states and our consideration shows that the resonant laser field can strongly affect all aspects of the impact ionization, including the angular and energy distributions of the emitted electrons, the total cross section and the projectile scattering.

Similar effects are expected to take place also in the impact ionization of more complex targets if the latter ones are affected by a resonant laser field. Such targets, like e.g. sodium atoms, can be better candidates for testing effects of a resonance laser field on the impact ionization in experiments on ion-atom collisions.

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