Laser-assisted and laser-induced electron-impact ionization during nonsequential double ionization of He

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(Received 13 September 2002; published 21 March 2003)

Two-electron sum-momentum distributions from nonsequential double ionization of He in an intense laser field are analyzed. The results of a *S*-matrix calculation show that the inelastic scattering of the first electron with the residual ion is best considered as a (internal) laser-induced impact ionization. The parallel component of the sum momentum is found to be largest when the prescattering electron energy is close to zero. This is due to the fact that a large momentum transfer between the sum momentum in the final state and the impact momentum in the intermediate state along the polarization axis favors the absorption of field energy by the two-electron system during the scattering process.

DOI: 10.1103/PhysRevA.67.035401 PACS number(s): 32.80.Rm, 32.80.Wr, 34.50.Rk, 34.80.Qb

Nonsequential double ionization of atoms is currently one of the most vigorously discussed problems in strong-field atomic physics. Recently, numerical results, based on the correlated energy sharing mechanism [1–3], for the recoil-ion momentum as well as electron momenta and energy distributions are found [2,3] to be in good agreement with experimental observations [4–6]. The correlated energy sharing mechanism is visualized for the case of double ionization of the He atom by the Feynman diagram in Fig. 1: first one of the two electrons becomes active and propagates after the absorption of photon energy from the field at time t_1 in the intermediate Volkov or field dressed states $\{\boldsymbol{k}\}$ [7], while the second electron propagates in the virtual states of the ion $\{j\}$. Then, the active electron shares its energy via the electron-electron correlation interaction $V_C = 1/r_{12}$ with the other electron at time t_2, resulting in the escape of both electrons together with momenta $\boldsymbol{k}_a$ and $\boldsymbol{k}_b$. In the final state of the process the electrons are affected by the laser field and are best described by Volkov states [2]. The diagram contains [2,8] both the antenna model [9] and the rescattering mechanism [10].

Note that the second step of the double ionization mechanism, distinguished by the shadowed part in Fig. 1, essentially represents an (internal) electron-impact ionization of the residual ion in the presence of a strong laser field at different incident energies. The distribution of the incident energies is determined by the amplitudes for absorption of different numbers of photons from the field by the active electron in the (internal) above-threshold-ionization-(ATI-)like process in the first step of the process. If the drift energy of the active electron $k^2/2$ exceeds the energy needed to liberate the second electron, we may speak of a laser-assisted ionizational collision, otherwise of a laser-induced impact ionization.

In this paper, we investigate the relative role of active electrons with different drift energies in the intermediate state on the double-ionization process by performing calculations for the sum-momentum distributions of the two electrons emitted from He. The amplitude, corresponding to the Feynman diagram in Fig. 1, can be written as [2,3] (Hartree atomic units, $\hbar = e = m = 1$, are used)

$$\frac{dW_N^{(NSDI)}(\boldsymbol{k}_a,\boldsymbol{k}_b)}{d\boldsymbol{k}_a d\boldsymbol{k}_b} = 2\pi\delta\left(\frac{k_a^2}{2}+\frac{k_b^2}{2}+E_{1S}+2U_p-N\omega\right) \times |T_N^{(NSDI)}(\boldsymbol{k}_a,\boldsymbol{k}_b)|^2, \tag{1}$$

where $T_N^{(NSDI)}$ is given by

$$T_N^{(NSDI)}(\boldsymbol{k}_a,\boldsymbol{k}_b) = \sum_n \sum_j \int d\boldsymbol{k}\frac{1}{\sqrt{2}}\langle \Phi^0(\boldsymbol{k}_a,\boldsymbol{r}_1)\Phi^0(\boldsymbol{k}_b,\boldsymbol{r}_2) + \boldsymbol{k}_a \leftrightarrow \boldsymbol{k}_b | V_C | \Phi_j^+(\boldsymbol{r}_2)\Phi^0(\boldsymbol{k},\boldsymbol{r}_1)\rangle (U_p - n\omega)$$

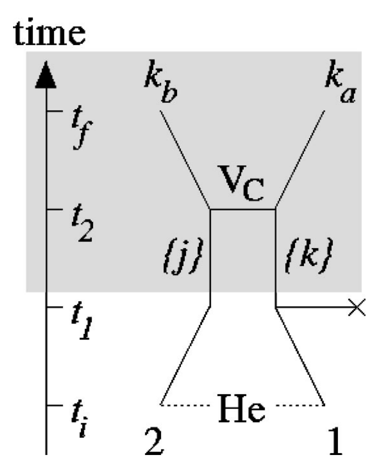

FIG. 1. Dominant Feynman diagram of nonsequential double ionization of He in an intense laser field; shadowed part indicates the (internal) $(e,2e)$-like process in the presence of the laser field.

$$\times \frac{J_{N-n}\left(\boldsymbol{\alpha}_0\cdot(\boldsymbol{k}_a+\boldsymbol{k}_b-\boldsymbol{k});\frac{U_p}{2\omega}\right)J_n\left(\boldsymbol{\alpha}_0\cdot\boldsymbol{k};\frac{U_p}{2\omega}\right)}{\frac{k^2}{2}-E_j+E_{1S}+U_p-n\omega+i0}$$
$$\times\langle\Phi_j^+(\boldsymbol{r}_2)\Phi^0(\boldsymbol{k},\boldsymbol{r}_1)|\Phi_{1S}(\boldsymbol{r}_1,\boldsymbol{r}_2)\rangle. \quad (2)$$

$\Phi^0(\boldsymbol{k},\boldsymbol{r})$ is the plane-wave state of momentum $\boldsymbol{k}$, $\Phi_j^+(\boldsymbol{r})$ are the intermediate states of the He^+ ion and $\Phi_{1S}(\boldsymbol{r}_1,\boldsymbol{r}_2)$ is the initial ground-state wave function of the He atom. $J_N(x,y)$ is the generalized Bessel function of two arguments, $U_p=I/4\omega^2$ denotes the ponderomotive energy (or quiver energy) of a free electron in a laser field of intensity I and frequency ω, while $\alpha_0=\sqrt{I}/\omega$ is its quiver radius. $E_{1S}=79.02$ eV and E_j are the ground-state binding energy of He and the binding energy of the intermediate He^+ state, respectively, and $V_C=1/r_{12}$ is the electron-electron correlation interaction.

For the actual calculations analytic atomic wave functions [11] are used, the sixfold space integrations in Eq. (2) are carried out analytically, the radial integration in k is performed by pole approximation, and the two integrals over the angles of $\boldsymbol{k}$ as well as the sum over n are performed numerically. The contribution of the lowest term of the sum over j, corresponding to the ground state of the He^+ ion, is only retained since it is found to dominate over the contributions from any excited state [1].

For the analysis of the results obtained for the momentum distributions from nonsequential double ionization, we found it instructive to investigate the (sub)process of electron-impact ionization of the ion in a laser field itself as well. The latter is a basic strong-field process (for a recent review, see, e.g., Ref. [12]) and elaborate numerical studies of the role of laser dressing of the target states on the sixfold differential cross sections have been performed (e.g., [13]). Using these theories to obtain the momentum distributions of present interest would result in time-consuming calculations. Since here, we are interested in a qualitative physical understanding of the relative role of laser-assisted and laser-induced impact ionization during the more complex double ionization process, we have restricted ourselves to perform calculations using a first Born approximation treatment for the scattering process and Volkov plane-wave states for the incident and the ejected electrons. The rate of electron impact ionization of ground state He^+ with initial electron momentum $\boldsymbol{k}_0$ is then given as

$$\frac{dW_N^{(e\text{-}2e)}(\boldsymbol{k}_a,\boldsymbol{k}_b,\boldsymbol{k}_0)}{d\boldsymbol{k}_a d\boldsymbol{k}_b}=2\pi\delta\left(\frac{k_a^2}{2}+\frac{k_b^2}{2}+E_{1s}^+ +U_p-\frac{k_0^2}{2}-N\omega\right)$$
$$\times|T_N^{(e\text{-}2e)}(\boldsymbol{k}_a,\boldsymbol{k}_b,\boldsymbol{k}_0)|^2, \quad (3)$$

with

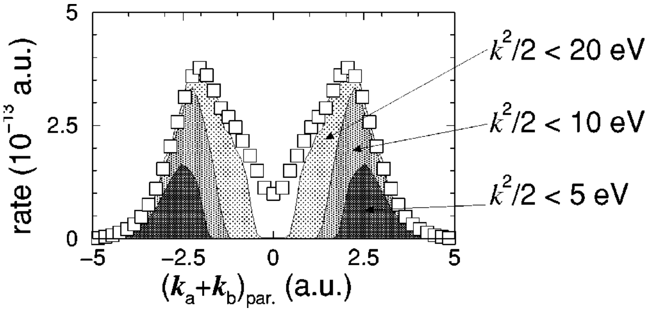

FIG. 2. Two-electron sum-momentum distributions from nonsequential double ionization of He atom parallel to the polarization axis of the field; laser parameters: $\lambda=800$ nm, $I=6.6\times10^{14}$ W/cm^2. Results of full S-matrix calculations (squares) are compared with those of restricted calculations, in which the value of the kinetic energy (and momentum k) of the Volkov electron in the intermediate states was limited.

$$T_N^{(e\text{-}2e)}(\boldsymbol{k}_a,\boldsymbol{k}_b,\boldsymbol{k}_0)=\frac{1}{\sqrt{2}}\Big\langle\Phi^0(\boldsymbol{k}_a,\boldsymbol{r}_1)\Phi^0(\boldsymbol{k}_b,\boldsymbol{r}_2)$$
$$+\boldsymbol{k}_a\leftrightarrow\boldsymbol{k}_b\left|\frac{1}{r_{12}}-\frac{2}{r_1}\right|\Phi_{1s}^+(\boldsymbol{r}_2)\Phi^0(\boldsymbol{k}_0,\boldsymbol{r}_1)\Big\rangle$$
$$\times J_N\left(\boldsymbol{\alpha}_0\cdot(\boldsymbol{k}_a+\boldsymbol{k}_b-\boldsymbol{k}_0);\frac{U_p}{2\omega}\right). \quad (4)$$

The two-electron sum-momentum distributions parallel and perpendicular to the laser polarization axis are obtained by integration of the sixfold differential rates, Eqs. (1) and (3), over the remaining momentum components and the summation over all N satisfying the δ constraints. The integrations are carried out using the Monte Carlo method.

We have performed calculations of the two-electron sum-momentum distributions for He atom parallel and perpendicular to the polarization axis of a linearly polarized Ti:sapphire laser field. In all calculations, we have chosen the laser parameters to be identical to those in the observations of the recoil-ion momentum from He atom by Weber *et al.* [4], i.e., $\lambda=800$ nm and $I=6.6\times10^{14}$ W/cm^2. Note that the results of the present S-matrix theory were found [2] to agree well with the experimental observations [4] at these laser parameters.

Here, our aim is to analyze the relative role of active electrons with different drift energies in the intermediate Volkov states, $\{\boldsymbol{k}\}$, on the nonsequential process and the experimental observables. To this end, we have calculated hypothetical momentum distributions by deliberately *restricting* the absolute value of the intermediate electron drift momentum k in Eq. (2) and compared them with those of the full calculations.

We first consider the two-electron sum-momentum distributions parallel to the polarization axis. The hypothetical distributions for three different upper limits of the intermediate electron energy, namely, $k^2/2<5$ eV, 10 eV, and 20 eV, are shown in Fig. 2 (hatched areas) along with those of the full calculation (squares). The restrictions have a drastic effect on the momentum distributions. It is seen from the figure that

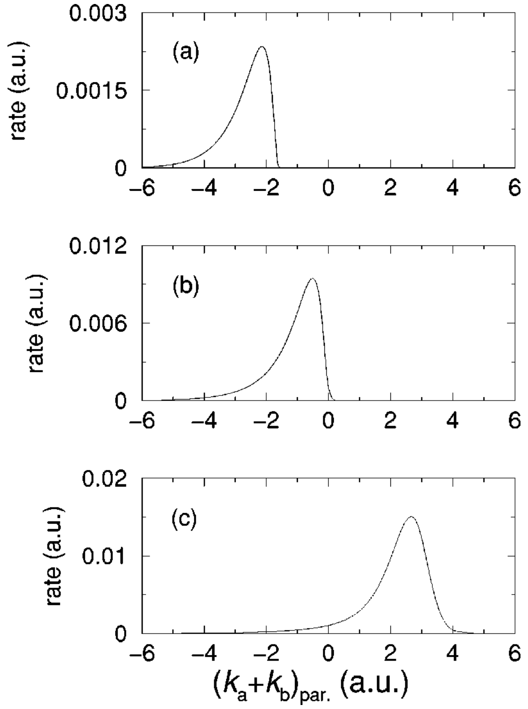

FIG. 3. Two-electron sum-momentum distributions from $(e,2e)$ processes in the direction parallel to the polarization axis. The incident electron momentum is chosen to be in the direction of the positive z axis (polarization axis). The panels correspond to kinetic energies of the incident electron equal to (a) 5 eV, (b) 30 eV, and (c) 200 eV. The laser parameters are $\lambda=800$ nm and $I=6.6\times10^{14}$ W/cm^2.

the parallel momentum distributions are "filled" from the outside to the center with increasing intermediate state (prescattering) electron energy. The component of the sum momentum parallel to the field is obviously largest when the drift energy of the active incident electron is close to zero.

For the present laser parameters, we retained the results of the full calculations (within the numerical accuracy of the present Monte Carlo calculations of about $\pm10\%$) as soon as we extended the energy limit to $k^2/2\approx50$ eV or more. From these results, we may conclude that the nonsequential double ionization process is dominated by (internal) scattering processes with impact energies of the active electron well below the ionization threshold of the He^+ ion. Thus, for the present laser parameters the second step of the nonsequential double ionization mechanism can be considered to be rather a laser-induced electron-impact ionization than a laser-assisted process.

To gain further insights into the above results, we have performed additional calculations of the sum-momentum distributions from electron-impact ionization of He^+ in a laser field. We have considered electron impact in the direction of the *positive* z axis (i.e., the polarization axis), since in the (internal) ATI-like ionization electrons are mainly ejected along the polarization axis (c.f., e.g., Ref. [4]). The results of our calculations for the parallel component of sum momentum are shown in Fig. 3; panels (a), (b), and (c) correspond to different energies of the incident electron, namely, 5 eV, 30 eV, and 200 eV, respectively.

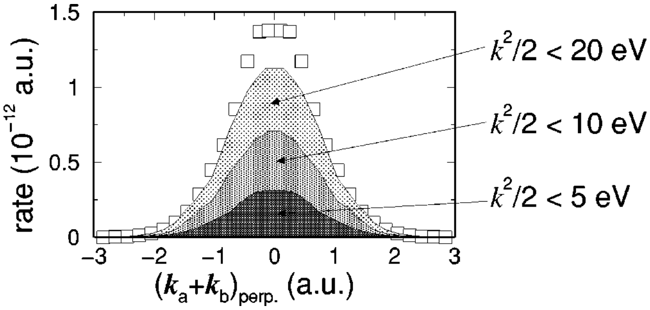

FIG. 4. Two-electron sum-momentum distributions from nonsequential double ionization of He atom perpendicular to the laser polarization axis; the rest as in Fig. 2.

One can see that in all cases the shapes of the distributions are narrow and rather similar. The position of the peak, however, changes largely with change of the kinetic energy of the incident electron. Whereas for $E_0=k_0^2/2$ equal to 5 eV [panel (a)] and 30 eV [panel (b)] the maximum occurs for a sum momentum, whose parallel component is in opposite direction to the momentum of the incident electron, for the highest electron energy [200 eV, panel (c)] the largest rate is found for a parallel component in the same direction as the incident electron momentum.

The shift of the maximum is closely related to the relation of the kinetic energy of the incident electron E_0 to the threshold energy $E_{tr}=E_{1s}+U_p=93.87$ eV needed to liberate the second electron. If the initial energy exceeds the threshold energy the maximum of the sum-momentum distribution appears at $(\boldsymbol{k}_a+\boldsymbol{k}_b)_{par.}\approx\sqrt{2(E_0-E_{tr})}=2.79$ a.u. [for $E_0=200$ eV, Fig. 3(c)]. This indicates a high probability for an exchange of a net number of zero photons between the field and the two-electron system, as expected for a laser-assisted process.

On the contrary, for small initial energies with $E_0<E_{tr}$ the two electrons need to absorb photons from the field to escape from the residual ion. The probability to achieve the required energy transfer is largest for large momentum transfers along the polarization direction between the two-electron sum momentum $\boldsymbol{k}_a+\boldsymbol{k}_b$ and the initial momentum $\boldsymbol{k}_0$. This can be readily understood from the cut off property of the generalized Bessel function in Eq. (4), that it tends to have significant values only up to orders that are comparable to their first argument and fall-off exponentially beyond [14]. Thus, rate contributions for the absorption of a large number of photons N from the field are most relevant for large arguments $\boldsymbol{\alpha}_0\cdot(\boldsymbol{k}_a+\boldsymbol{k}_b-\boldsymbol{k}_0)$, or, equivalently in the present case of electron impact along the polarization direction, when $(\boldsymbol{k}_a+\boldsymbol{k}_b)\|-\boldsymbol{k}_0$. Further, the absolute value of the sum momentum is expected to be larger, the smaller the initial kinetic energy is and, consequently, the larger has to be the support from the field to initiate the $(e,2e)$ process, which is what we observe from the results in Fig. 3(a,b). These conclusions are in agreement with the results for the distributions for nonsequential double ionization (c.f., Fig. 2), that a restriction of the intermediate-state electron energy to low values lead to large sum-momentum components along the polarization axis. Our results concerning the kinematics

also agree with considerations based on adiabatic transition theory [9(c)].

We further note that the above findings are in close analogy to previous results from laser-assisted elastic-scattering processes, in which an electron, incident along the polarization axis, gains maximum energy from the field when it is backscattered [16]. In this respect, it is the sum momentum of the two electrons that replaces in the nonsequential double ionization the momentum of the scattered electron in elastic rescattering of ATI.

Finally, we present in Fig. 4 the results of our calculations for the perpendicular component of the sum-momentum distribution from nonsequential double ionization. The restriction of the kinetic energy in the intermediate state does not affect the overall shape of the distribution, but only reduces its absolute rate. This may be expected, since the field does not couple to the electron momentum components perpendicular to the polarization axis [2] and, hence, the strength of the field support during the impact ionization process does not affect the perpendicular momentum distributions.

In conclusion, results of numerical calculations for the two-electron sum-momentum distributions show that laser-induced processes with incident electrons having drift energies well below the ionization threshold of He^+ ion dominate the nonsequential double ionization process. It is further seen that small prescattering electron energies in the intermediate state lead to a large sum momentum of the two electrons in the final state. This is due to the fact that the required additional energy transfer from the field to liberate the second electron is most probable, when the final-state two-electron sum-momentum is large and antiparallel to the electron momentum before the scattering process.

This work has been supported by the Alexander von Humboldt Foundation, the Polish Committee of Science, Grant No. KBN 2P03B03919, and the Deutsche Forschungsgemeinschaft, Grant No. BE 2092/6-1.

[1] A. Becker and F.H.M. Faisal, J. Phys. B **29**, L197 (1996).
[2] A. Becker and F.H.M. Faisal, Phys. Rev. Lett. **84**, 3546 (2000).
[3] A. Becker and F.H.M. Faisal, Phys. Rev. Lett. **89**, 193003 (2002).
[4] Th. Weber *et al.*, Phys. Rev. Lett. **84**, 443 (2000).
[5] Th. Weber *et al.*, Nature (London) **405**, 658 (2000).
[6] R. Lafon *et al.*, Phys. Rev. Lett. **86**, 2762 (2001).
[7] D.M. Volkov, Z. Phys. **9**, 250 (1935).
[8] A. Becker and F.H.M. Faisal, Phys. Rev. A **59**, R1742 (1999).
[9] M.Yu. Kuchiev, Pis'ma Zh. Eksp. Teor. Fiz. **45** 319 (1987) [JETP Lett. **45**, 404 (1987)]; J. Phys. B **28**, 5093 (1995); Phys. Lett. A **212**, 77 (1996).
[10] P.B. Corkum, Phys. Rev. Lett. **71**, 1994 (1993).
[11] E. Clementi and C. Roetti, At. Data Nucl. Data Tables **14**, 177 (1974).
[12] F. Ehlotzky, A. Jaroń, and J.Z. Kamiński, Phys. Rep. **297**, 63 (1998).
[13] C.J. Joachain *et al.*, Phys. Rev. Lett. **61**, 165 (1988); O. El Akramine *et al.*, J. Phys. B **32**, 2783 (1999); A. Makhoute *et al.*, *ibid.* **32**, 3255 (1999).
[14] This cut-off property is well known (e.g., Ref. [15], p. 87) for the ordinary Bessel function and is found [2] to hold true for the generalized Bessel functions in the regime of field parameters of present interest as well.
[15] *Higher Transcendental Functions*, edited by A. Erdelyi (McGraw-Hill, New York, 1953), Vol. 2.
[16] G.G. Paulus *et al.*, J. Phys. B **27**, L703 (1994).