Laser assisted (e, 2e) process of the helium atom

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Abstract. The dynamics of the electron impact multiphoton ionization of a Helium atom in the presence of an intense laser field \((n\gamma, 2e)\) is studied theoretically. Significant modifications are noted in the laser assisted cross sections with respect to the field free ones. The results are compared with the recent kinematically complete experiments for high incident energies (1 KeV).

1. Introduction

Laser assisted (LA) excitation and ionization of atoms by charged particles plays a very dominant role in many applied areas such as fusion plasma, plasma heating, high power gas lasers, semiconductor physics etc [1]. Further, the LA electron-atom collision allows experimental observations [2-4] of different multiphoton (MP) phenomenon at relatively moderate field strengths. Despite such immense importance, the detailed dynamics of such LA processes is far from being well understood and is now a subject of great challenge.

A number of theoretical works have been put forward for the LA (e, 2e) process of the H and He atoms both prior [5-7] to and following [8-10] the recent kinematically complete experiment of Hörr et al [2].

2. Theory

The following laser assisted (LA) multiphoton ionization process of Helium atom in its ground state is studied for exchange of \(n\) photons with frequency \(\omega\) and field strength \(E_0\):

\[
e^-(E_1, k_1) + He^- (1s) \pm n\gamma (\omega, \vec{E}_0) \rightarrow e^-(E_1, \vec{k}_1) + e^+(E_2, \vec{k}_2) + He^*
\]

(1)

The laser field is chosen to be monochromatic, spatially homogeneous, linearly polarized and is represented by \(\vec{E}(t) = \vec{E}_0 \sin (\omega t + \delta)\) corresponding to the vector potential in the Coulomb gauge \(\vec{A}(t) = \vec{A}_0 \cos (\omega t + \delta)\) with \(\vec{A}_0 = c \vec{E}_0 / \omega\) and \(\delta\) is the initial phase of the laser field. The laser polarization is chosen to be parallel to the projectile momentum.

The prior form of the transition matrix element for the process (1) is given as[10]:

\[
T_\gamma = -i \int_{-\infty}^{\infty} dt \langle \Psi_j^- | V_i | \psi_i \rangle
\]

(2)

where \(V_i\) is the perturbation potential in the initial channel.

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$V_i = -\frac{2}{r_1} + \frac{1}{r_{12}} + \frac{1}{r_{13}}$ \hfill (3)

$r_1$, $r_2$, $r_3$ are the position vectors of the projectile electron and the two target electrons (2, 3) respectively, $r_{12} = r_1 - r_2$ and $\vec{r}_{13} = \vec{r}_1 - \vec{r}_3$, $\psi_i$ is the asymptotic initial state wavefunction and $\Psi_f$ is the final state wavefunction satisfying the incoming wave boundary condition. It is evident from Eq. (2) that the perturbation $V_i$ vanishes asymptotically for $r_1 \to 0$, $r_2, r_3$ finite. In the present model the role of the external laser field is to modify the projectile – electron as well as the ground and residual target atom / ion wavefunctions.

Certain distinctive features of the present model are as follows:

(1) Projectile residual target nucleus interaction in the final channel is chosen to be the Coulomb Volkov (CV) wavefunction [7]. The approximate CV solution chosen for the scattered and ejected electron presents the advantage of containing the fields in all orders (incoherent in the plane wave Volkov solution), although it completely decouples the projectile-laser and the projectile-screened ion interaction. For low laser intensity or high laser frequency, the coupling of the slow electron with the laser field is much weaker than the Coulomb interaction so that a perturbative treatment can be used.

(2) Target dressing is taken into account both in the initial and final channel in the framework of first order time dependent perturbation theory, which is quite legitimate since the present laser field strength $\mathcal{E}_0 \ll$ a.u. of field strength ($5 \times 10^{11}$ V/m).

(3) Apart from the simple CV, a more refined wavefunction, the Modified Coulomb Volkov (MCV) is considered [6,10] for the dressed ejected electron. For low energy ejected electron and/or for strong laser field with low frequency the CV wavefunction is not adequate, a modification is needed [6].

(4) Correlation between the scattered and the ejected electrons is taken into account in the final channel so that the asymptotic three-body boundary condition is satisfied which is essential for a reliable estimate of an ionization process.

The expression for the triple differential cross-section (TDCS) for n photon transfer is given as:

$$\frac{d^3\sigma}{d\Omega_1 d\Omega_2 dE_3} = 2 \frac{k_1 k_2}{k_i} | T^n_{\theta} |^2$$

3. Results & Discussions

The triple differential cross-sections (TDCS) are computed for the process (1) with the kinematics chosen mainly in accordance with the experiment [2].

**Figure 1.** TDCS for $l=1$ as a function of ejected electron angle ($\theta_2$). The incident energy is $E_i = 1000$ eV, the ejected electron energy $E_2 = 3.7$ eV and the scattering angle is $\theta_i = 6.4^\circ$. Field strength $\mathcal{E}_0 = 5\times 10^9$ V/m, laser photon energy $\omega = 1.17$ eV. Dashed line: FF results, dotted line: CV results and solid line: MCV results.
Figure 1 reveals that for single photon transfer (\(l = 1\)), the CV results are strongly suppressed compared to the FF ones. For the MCV case, on the other hand, the binary peak is suppressed while the recoil peak is enhanced with respect to the FF results.

Figure 2 compares the CV multiphoton TDCS (MTDCS) with the experiment [2], where the present binary peak is normalized to the experimental one. The present results qualitatively follows the experimental behaviour except in the recoil region although the absolute value of the present results overestimates the experiment both in the binary and the recoil region.

Figure 3 exhibits the behaviour of the MTDCS w.r.t. the laser frequency (\(\omega\)) at \(E_i = 1000 \text{ eV}\) and it reveals that the magnitude of the MTDCS decreases with increasing laser frequencies as is expected physically. For higher laser frequency, the electron-laser coupling parameter \(\alpha_0 = \frac{e_0}{\omega^2}\) decreases resulting in lesser laser modifications of the cross-section from the field free values. Since the MCV results are much enhanced compared to the FF ones, the former decreases with increasing laser frequency.

Mathematically, the above feature could also be explained from the properties of the Bessel functions [11, 12] occurring in the transition amplitude \(T_{\theta_i}\). The argument of the Bessel functions \(\left(k_i \cdot \tilde{k}_j - \vec{\alpha}_0\right)\) decreases for higher laser frequency i.e. for lower value of \(\alpha_0\). Hence, with decreasing value of the argument \(J_0(x)\) increases while \(J_{\ell}(x)\) (for \(\ell \neq 0\)) decreases.
so that when the individual contributions due to different $\ell$ values (e.g. $0, \pm 1, \pm 2, \pm 3$) are added in the summed TDCS (MTDCS) the net effect is that the summed cross-section diminishes.

Further, the wavefunction of the ejected electron occurring in $\psi_f$ of Eq. (2) contains a term $k_z \tilde{\alpha}_0$ (vide Eq.5 of [10]). This also explains the above behaviour of the TDCS w.r.t frequency.

For low frequency e.g. $\omega = 1.17$ eV ($\alpha_0 \sim 5.25$), there is significant deviation between the CV and the MCV results while the deviation diminishes with increasing frequency / decreasing $\alpha_0$ (electron–laser coupling) (vide figure 3) so that for $\alpha_0 \sim 0.64$ the two (CV & MCV) almost coincide.

4. Conclusions:
A strong enhancement is noted in the present multiphoton binary peak intensity, in qualitative agreement with the experiment. The deviation between the present CV and MCV multiphoton TDCS becomes negligible for higher frequency / lower field strength (i.e. $\alpha_0 < 1$).

5. References:
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