Coherent phase control in electron–argon scattering in a bichromatic laser field*

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This paper theoretically investigates the coherent phase control in electron–argon scattering assisted by a bichromatic laser field. The laser field is composed of a fundamental component and its second harmonic. The incoming and out going states of electron are described by the Volkov wave functions, and the electron–target interaction is treated as a screening potential. Numerical results for differential cross section of multiphoton processes vs the phase difference between the two components of laser field are discussed for several scattering angles and impact energies.

Keywords: coherent phase control, bichromatic laser field, electron–argon, free–free transition
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Multiphoton free–free transitions (MFFT) have attracted much attention since the pioneering work by Bunkin and Fedorov[1] and Kroll and Watson,[2] and a great deal of papers has been devoted to them.[3−9] As the experimental technology improved, powerful lasers and new kinds of laser fields have been applied to laser-assisted atomic and molecular processes. Due to the significance in dynamic control of the multicolour lasers, the atomic and molecular processes in them received considerable attentions.[10−20] One can enhance or modify the processes by changing the phase difference between the laser components. In most of the theoretical work, the laser radiation is treated as a classical electromagnetic field. This relies on the argument that in a laser beam the density of radiation quanta is so large that the depletion of this beam by emitting or absorbing quanta from it is negligible.[21]

In this paper, we consider the coherent phase control in the free–free transition (FFT) in argon atom assisted by a bichromatic laser field. The laser field is treated classically as a electromagnetic field of two components of frequencies, $\boldsymbol{\mathcal{E}}(t) = \boldsymbol{\mathcal{E}}_0[\sin(\omega t) + \sin(2\omega t + \varphi)]$, where $\boldsymbol{\mathcal{E}}_0$ is the electric field amplitude vector and the relative phase φ can be arbitrarily changed. Atomic units $\hbar = m = e = 1$ are used throughout.

We suppose that the laser frequency and intensity are sufficiently low so that the excitations of atomic transitions can be neglected, and the atomic target can be described by a screening potential[8,22,23]

$$V(r) = -\frac{Z}{r}\sum_{i=1}^{3} A_i \exp(-\alpha_i r), \tag{1}$$

where r denotes the position of the electron with respect to the nucleus, and Z denotes the nuclear charge number. For argon, $A_1 = 2.1912$, $A_2 = -2.8252$, $A_3 = 1 - A_1 - A_2$, $\alpha_1 = 5.5470$, $\alpha_2 = 4.5687$, and $\alpha_3 = 2.0446$.

The scattering matrix for the laser-assisted FFT in the first Born approximation reads as

$$S_{\mathrm{f\,i}}^{(1)} = -\mathrm{i}\int_{-\infty}^{\infty} \mathrm{d}t \int \mathrm{d}\boldsymbol{r}\chi_{\boldsymbol{k}_{\mathrm{f}}}^{*}(\boldsymbol{r},t)V(r)\chi_{\boldsymbol{k}_{\mathrm{i}}}(\boldsymbol{r},t), \tag{2}$$

where $\chi_{\boldsymbol{k}_{\mathrm{f}}}$ and $\chi_{\boldsymbol{k}_{\mathrm{i}}}$ are the initial and final states of electron respectively, which are described by the Volkov wave function

$$\begin{aligned}&\chi_{\boldsymbol{k}_{\mathrm{i,f}}}(\boldsymbol{r},t)\\ &= \exp\left\{\mathrm{i}\left[\boldsymbol{k}_{\mathrm{i,f}}\cdot\boldsymbol{r} - E_{\mathrm{i,f}}t - \frac{1}{\omega^2}\boldsymbol{k}_{\mathrm{i,f}}\cdot\boldsymbol{\mathcal{E}}_0\sin(\omega t)\right.\right.\\ &\quad\left.\left. -\frac{1}{4\omega^2}\boldsymbol{k}_{\mathrm{i,f}}\cdot\boldsymbol{\mathcal{E}}_0\sin(2\omega t+\varphi)\right]\right\},\end{aligned} \tag{3}$$

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where $\boldsymbol{k}_{\rm i,f}$ are the wave vectors of electron in the initial and final states, and the electron energies $E_{\rm i,f} = k_{\rm i,f}^2/2$.

Using the potential of Eq. (1) and the wave functions of Eq. (3), we obtain

$$S_{\rm fi}^{(1)} = -2\pi {\rm i} \sum_l T_{\rm f,i}^{(1)}(l)\delta(E_{\rm f} - E_{\rm i} + l\omega). \qquad (4)$$

Where $T_{\rm f,i}^{(1)}(l)$ is the transition matrix element, resolved with respect to multiphoton exchange processes

$$T_{\rm f,i}^{(1)}(l) = {\rm B}_l\left(\lambda, \frac{1}{4}\lambda, \varphi\right) V(k_{\rm f,i}), \qquad (5)$$

in which

$$V(k_{\rm f,i}) = \int {\rm d}\boldsymbol{r}\, {\rm e}^{-{\rm i}(\boldsymbol{k}_{\rm f} - \boldsymbol{k}_{\rm i})\cdot \boldsymbol{r}} V(\boldsymbol{r}) \qquad (6)$$

is the Fourier transformation of the potential, and

$$\begin{aligned} &{\rm B}_l\left(\lambda, \frac{1}{4}\lambda, \varphi\right) \\ &= \sum_{n=-\infty}^{\infty} {\rm J}_{l-2n}(\lambda){\rm J}_n\left(\frac{1}{4}\lambda\right)\exp(-{\rm i}n\varphi) \end{aligned} \qquad (7)$$

is the generalized Bessel function with $\lambda = (\boldsymbol{k}_{\rm f} - \boldsymbol{k}_{\rm i}) \cdot \boldsymbol{\mathcal{E}}_0/\omega^2$. The ${\rm J}_n(\lambda)$ is the ordinary Bessel function. The differential cross section (DCS) for the net exchange of $l\omega$ energy between the scattering system and the bichromatic laser field reads as

$$\frac{{\rm d}\sigma}{{\rm d}\Omega} = \frac{1}{(2\pi)^2}\frac{k_{\rm f}}{k_{\rm i}}|T_{\rm f,i}^{(1)}(l)|^2. \qquad (8)$$

Equation (7) implies that the transition amplitude for $l\omega$ energy exchanged consists of many terms. Each term represents a process with the absorption of n harmonic photons and $l - 2n$ fundamental component photons. These terms have interference through the phase factors $\exp(-{\rm i}n\varphi)$. Therefore the DCS of Eq. (8) depends on the relative phase φ between two laser components.

For numerical calculation, we study the electron–argon scattering under the geometry of the experiment of Weingartshofer *et al.*[5] in which the angle between the polarization vector $\boldsymbol{\mathcal{E}}_0$ and the momentum $\boldsymbol{p}_{\rm i}$ of the incident electron is $\psi_0 = 38°$, the momentum $\boldsymbol{p}_{\rm f}$ of the scattered electron is in the plane defined by the polarization vector $\boldsymbol{\mathcal{E}}_0$ and the momentum $\boldsymbol{p}_{\rm i}$ of the incident electron. We choose the field frequency $\omega = 0.117$ eV (CO_2 laser) and the amplitude $\boldsymbol{\mathcal{E}}_0 = 2.7 \times 10^5\ {\rm V}\cdot{\rm cm}^{-1}$.

In Fig. 1, we display the relative phase dependence of DCS for each multiphoton process. At the field strength considered, only the processes with a few exchanged photon make significant contributions. With increasing $|l|$, the DCS declines rapidly. The variation tendencies for $l > 0$ (photon emission or bremsstrahlung) and $l < 0$ (photon absorption or inverse bremsstrahlung) are opposite. The results for bremsstrahlung attain the maxima at $\varphi = 180°$, while those for inverse bremsstrahlung attain their minima. This may attribute to the fact that the generalized Bessel function satisfies the relation:[10]

$${\rm B}_{-l}(a, b, \varphi) = (-1)^l {\rm B}_l^*(a, b, \varphi - \pi). \qquad (9)$$

In addition, the results for $l = 0$ and odd numbers are less sensitive to the variation of φ than $l =$ even. Generally speaking, the value of l roughly determines the terms to be taken into account in the sum of Eq. (7). For $l = 0$, the main contribution is the $n = 0$ term, therefore we observe no phase dependence in this case. For $l = 1$ and 3, the main contributions are respectively the terms of $n = 0, 1$ and $n = 0, 1, 2$. For $l = 2$ and 4, the main contributions are respectively $n = 0, 1$ and $n = 1, 2$. The relatively small number of terms and strong interference in the latter case make the φ-dependences more notable.

Figure 2 shows the results for $l = \pm 2$ at the scattering angles $\theta = 8°$, $13°$, $30°$, and $50°$. The DCS at $\theta = 30°$ is much more sensitive to φ than at other angles. At this angle, the result for bremsstrahlung is maximally enhanced at $\varphi = 180°$, while that for inverse bremsstrahlung is seriously depressed. In this case, one can effectively control the dynamics by varying the relative phase.

Figure 3 displays the result of $l = \pm 2$ for different impact energies at the scattering angle $\theta = 13°$ where the phase dependence is most sensitive. As we have observed in Fig. 1, the variation tendencies for bremsstrahlung ($l = +2$) and inverse bremsstrahlung ($l = -2$) are opposite. For the higher impact energy, the result is more sensitive to the variation of phase. When the impact energies increase, the scattered electron may exchange more photons with the laser field, while the results for $l = \pm 2$ only make a small contribution.

In summary, we have studied the electron–argon elastic scattering in a bichromatic laser field by employing a simple screening potential. The dependence of DCS on the relative phase φ is investigated at different scattering angles and impact energies. These are important in the coherent control in the FFT.

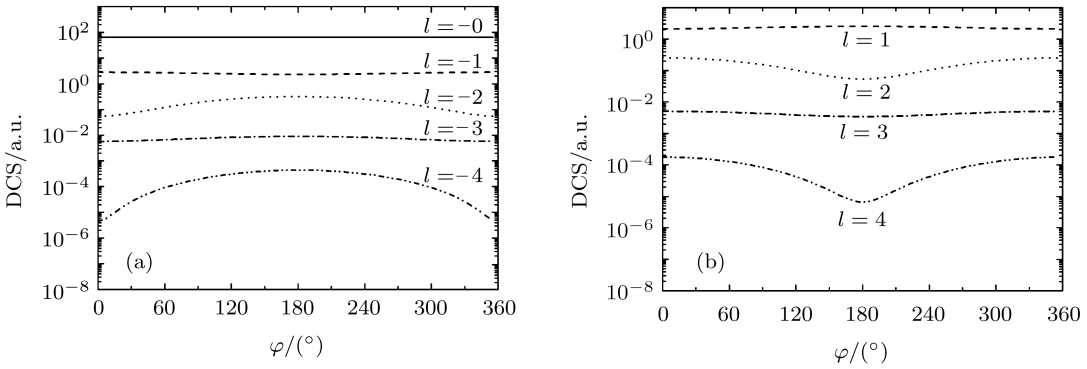


Fig. 1. The φ-dependence of multiphoton DCS at a scatteing angle $\theta = 13^\circ$ for electron–argon elastic scattering in a CO_2 laser field ($\hbar\omega = 0.117$ eV). The impact energy of the incident electron is $E_i = 9.5$ eV, and the laser amplitude $\boldsymbol{\mathcal{E}}_0 = 2.7 \times 10^5$ V·cm^{-1}. (a) results for bremsstrahlung ($l \geqslant 0$). (b) results for inverse bremsstrahlung ($l < 0$).

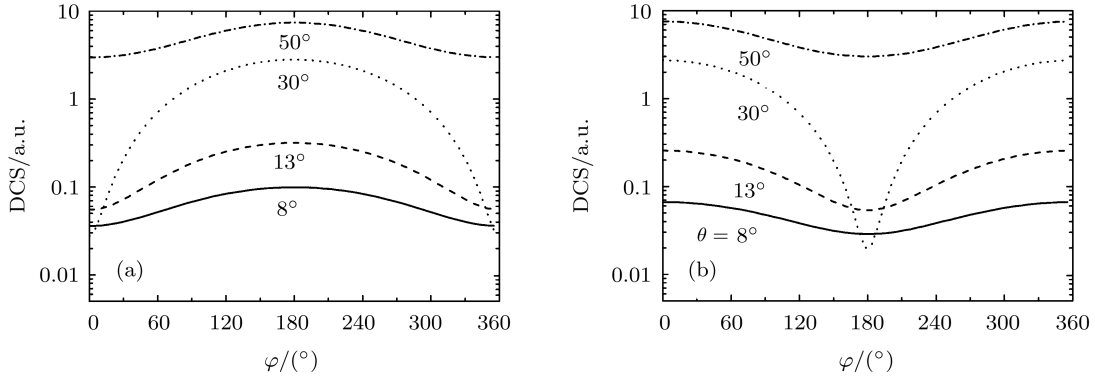


Fig. 2. The φ-dependence of DCS for $l = \pm 2$ at different scattering angles. The impact energy and laser parameters are the same as in Fig. 1. (a) results for bremsstrahlung ($l = 2$). (b) results for inverse bremsstrahlung ($l = -2$).

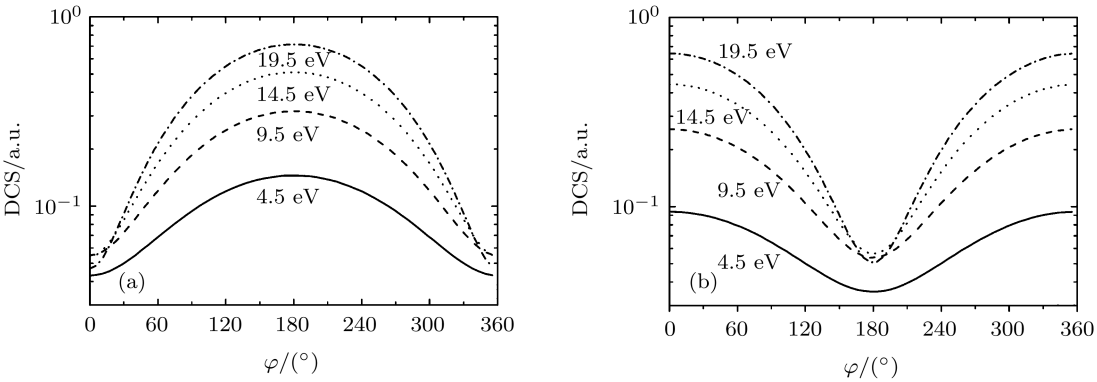


Fig. 3. The φ-dependence of DCS for $l = \pm 2$ at different impact energies. The scattering angle is $\theta = 13^\circ$. The laser parameters are the same as in Fig. 1. (a) results for bremsstrahlung ($l = 2$). (b) results for inverse bremsstrahlung ($l = -2$).

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