Production of intense isolated attosecond pulses with circular polarization by using counter-propagating relativistic lasers

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Abstract
We demonstrate theoretically and numerically that intense isolated circularly polarized (CP) attosecond pulses can be generated from ultrathin foil targets irradiated by two relativistic lasers from opposite sides, where their polarizations are orthogonal to each other. With a proper matching condition, the compressed oscillating plasma mirrors on both sides of the foil are pushed inside by laser radiation pressures, eventually merging together to form a dense electron nanobunch under the effect of orthogonal laser fields. This nanobunch reaches both high density and high energy in only half a laser cycle and smears out in others, resulting in coherent synchrotron emission of a single attosecond pulse with circular polarization. Two-dimensional particle-in-cell simulations show that an intense isolated CP attosecond XUV pulse with an intensity of $1.2 \times 10^{19}$ W cm$^{-2}$ and a duration of $\sim 75$ as can be obtained by two lasers with the same intensity of $2.1 \times 10^{20}$ W cm$^{-2}$.

1. Introduction

Controlling electrons with ultrashort intense laser pulses on the electron’s natural time scale, the attosecond ($1 \text{ as} = 10^{-18}$ s), has motivated the birth of a new science, attosecond physics [1–3]. To date, most investigations focus on only linearly polarized attosecond pulses obtained from relativistically oscillating mirror (ROM) mechanism [8], oblique irradiation on a solid target with intense CP or elliptically-polarized lasers [23, 24] are proposed to obtain a CP attosecond pulse, where the polarization of HHG is inherited from the drive laser. Previously, we have also proposed [25] to use an S-polarized two-color laser with a proper match of the ratio of the dichromatic driving field in different harmonics for production of CP attosecond pulses by ROM. However, for either CP or S-polarized lasers, the HHG conversion efficiency drops significantly. Furthermore, a so-called slingshot mechanism [26] has
been proposed recently to generate isolated CP attosecond pulses as the attosecond radiations always have come in the form of PHz-repetition-rate trains for the multi-cycle nature of the drive laser [27, 28].

Different from ROM, a more efficient HHG mechanism, coherent synchrotron emission (CSE) [29–32], has been identified in which a dense electron nanobunch with a δ-like peak density distribution and relativistic energy is formed, resulting in highly efficient radiation of XUV/x-rays. In CSE, the reflected radiation is proportional to the time derivative of the transverse current density in the nanobunch, instead of a simple phase modulation of the incident laser from plasma mirror in ROM. The peak reflected field can exceed that of the incident laser, resulting in a substantial increase in both radiation energy and harmonic orders, where the harmonic spectrum is characterized by a much slower decay scaling as $I(n) \propto n^{-8/3}$ for ROM. Besides, the interaction between the relativistic laser and the ultrathin foil to obtain HHG has been studied extensively [26, 33–35] with extra mechanism proposed including the sliding mirror mechanism [33] and the relativistic electron mirror mechanism [34]. However, the CSE mechanism is still the most efficient mechanism without the limitation on the longitudinal driving forces and the intensity of lasers.

In this paper, by theoretical modeling and two-dimensional (2D) particle-in-cell (PIC) simulations, we demonstrate, for the first time, that intense isolated circularly polarized (CP) attosecond pulses can be generated from ultrathin foil targets by irradiation from opposite sides with two relativistic lasers of orthogonal polarizations, where CSE dominates. In this scheme, two relativistically oscillating plasma mirrors formed on both sides of the foil at early stage are compressed and pushed inside by laser radiation pressures. With a proper matching condition between lasers and targets, the two plasma mirrors eventually merge together at the peak laser intensity cycle, forming a dense electron nanobunch under the effects of both laser fields with their polarization orthogonal to each other. CSE occurs when the nanobunch has the maximum compressed electron density $n_e$, meanwhile moving with the maximum longitudinal velocity $v_x$, which can be represented by a longitudinal Lorentz factor $\gamma_x = (1 - v_x^2/c^2)^{-1/2}$ at the same laser cycle. Afterwards, the foil target becomes transparent, leading to heating and decompression of the nanobunch. These make CSE substantially less efficient at later cycles, resulting in the production of a single (or few) bright CP attosecond pulse. 2D PIC simulations show that an intense isolated CP attosecond XUV pulse with an intensity of $1.2 \times 10^{20} \text{ W cm}^{-2}$ and a duration of $\sim 75\text{as}$ can be obtained by two lasers with the same intensity of $2.1 \times 10^{20} \text{ W cm}^{-2}$. Compared with the slingshot mechanism [26], the isolation of the CP attosecond in our scheme is controlled through the thickness of the target without the requirement of few-cycle CP driving lasers and a shorter attosecond pulse is achieved.

2. Theoretical analysis

The physical process of our scheme for production of intense isolated CP attosecond pulses is composed of three stages. At the first stage, ROM occurs at both sides of the foil, resulting in HHG on both sides, where the laser oscillating ponderomotive force of respectively $F_{\text{left}} = |E_1|^2 \cos(2\omega_0 t - 2k_0 \Delta x + 2\phi_0)$ and $F_{\text{right}} = |E_1|^2 \cos(2\omega_0 t + 2k_0 \Delta x)$ play roles. Here, $E_1$, $\omega_0$, $k_0$, $\phi_0$ and $\Delta x$ are the electric field strength, angular frequency, wave vector, the initial phase difference of the two lasers and the fluctuation of the initial position of the thin foil, respectively. Meanwhile, the two plasma mirrors are also compressed and pushed inside the foil by the laser steady ponderomotive force, i.e., radiation pressure $\sim |E_0|^2$. Due to the normal incidence of both lasers, at this stage, the conversion efficiency of HHG on both sides is rather low.

Secondly, around the cycle when both lasers reach their peak intensities, the above two plasma mirrors eventually merge together, forming a single dense electron nanobunch with thickness close to the skin depth, which undergoes effects of both laser fields with their polarization orthogonal to each other. CSE occurs when the nanobunch has the maximum compressed electron density $n_e$, meanwhile moving with the maximum longitudinal velocity $v_x$ at this laser cycle. The transverse radiation field $E_{\gamma,t}$ from the transverse current $J_{\gamma,t}$ has the form of $E_{\gamma,t}(x, t) = -1/2 \int_{-\infty}^{\infty} dx' J_{\gamma,t}(x', t')$. Utilizing the retardation relation $x' \pm t' = x \pm t$ and the relation $J_{\gamma,t} = -n_e v_x = -n_e P_{\gamma,t}/\gamma = -n_e P_{\gamma,t}/\sqrt{1 + P_{\gamma,t}^2 + P_{\gamma,t}^2 + P_{\gamma,t}^2}$, the radiation field becomes:

$$
E_{\gamma,t} = \frac{1}{2} n_e \gamma \left(1 \pm v_x\gamma\right) \left(P_{\gamma,t} + P_{\gamma,\gamma}\right) \left(1 + P_{\gamma,t}^2 + P_{\gamma,t}^2 + P_{\gamma,t}^2\right)^{1/2},
$$

where ‘−’ is for the radiation propagating to the left side and ‘+’ is for the radiation propagating to the right side. For the interaction between the relativistic laser and the ultrathin foil, according to the
the electrostatic restoring force \([36, 37]\), and laser vector potential, respectively. Should less than the skin depth of the transverse momentum is achieved \([29, 30]\), the high-order harmonics are emitted at the initial transverse positions of the electron bunch. To make sure the electrons in the foil can experience both lasers, the spatiotemporal points of the combined bunch. Therefore, the transverse momentum has the form

\[
p_y = a_y - \alpha(y - y_0) \quad (2)
\]

Here, \(p_y, p_z\) and \(a_y, a_z\) mean the transverse momentum of the electron bunch and the amplitude of the incident laser vector potential, respectively. \(\alpha = 2\pi n_e e^2 d_0/m_e\omega_0 c\) and \(y_0, z_0\) are the normalized areal density and the initial transverse positions of the electron bunch. \(d_0\) is the thickness of the ultrathin foil. \(c\) and \(m_e\) are the velocity of light in vacuum and the mass of electron. For the one-dimensional (1D) model, the transverse charge separation is neglected, so the transverse momenta become

\[
p_y = p_z = a_y. \quad (3)
\]

Then, the radiation field can be described as:

\[
E_{y,z} = \frac{n_e \gamma (1 \pm v_x)(a_y + a_z)}{1 + a_y^2 + a_z^2}. \quad (4)
\]

Dimensionless quantities are used: \(n_e = n_e/n_c, v = v/c, E = E/E_0 = eE/(m_e\omega_0 c), B = eB/(m_e\omega_0), a = eA/m_e c, n_e = \omega_0^2 e^2 m_e/e^2\) is the plasma critical density. As the merged electron nanobunch undergoes effects by both lasers together, \(P_{\text{CP}}\) can be estimated as \(a_y + a_z\) with \(a_y = a_0 \sin(t' - x')\), \(a_z = a_0 \cos(t' + x')\), where \(a_0\) means the magnitude of the incident laser potential vector and \((x', t')\) means the spatiotemporal points of the combined bunch. Therefore, the transverse momentum has the form

\[
p_y, z = \sqrt{p_y^2 + p_z^2} = a_0 \sqrt{\left[1 - \sin(2\pi \gamma \sin(2\pi t))\right]}. \quad (4)
\]

Because the target oscillates around \(x' = 0\) and the minimum of the transverse momentum is achieved \([29, 30]\), the high-order harmonics are emitted at \(x' = 0\). According the equation (4), the radiation spectra in the orthogonal directions will carry the phase difference of the two incident lasers. Therefore, the CP attosecond radiation is obtained. Considering the conservation of the areal density, the foil thickness \(d = n_e d_0/n_c^2\) during the compression, where \(n_c^2\) is the varying density of electron bunch. To make sure the electrons in the foil can experience both lasers, \(d'\) should less than the skin depth \(l_s = \lambda/2\pi \sqrt{n_c}\). As \(d'_{\text{min}} = d_0 - \lambda/\pi S\), the initial foil thickness should in principle satisfy \(d_0 \geq l_s + \lambda/\pi S\), where \(\lambda/\pi S\) is obtained through the balance between the Lorentz force and the electrostatic restoring force \([36, 37]\), and \(S = n_e/a_0\) is the similar parameter \([38]\). On the other hand, the foil should avoid transparency at the initial time, so that \(d_0 \geq \lambda/\pi S\) should be satisfied. Therefore, after some derivations, we obtain that the condition for our scheme to work is

\[
\frac{\lambda}{\pi S} \leq d_0 \leq \frac{\lambda}{2\pi S} \left(1 + \sqrt{\frac{S}{d_0} + 1}\right). \quad (5)
\]

If the foil thickness is even larger than the condition (5), two plasma mirrors could not be merged together to form a single relativistic nanobunch, leading to failure of CSE occurrence. If it is too thinner than the
condition (5), the foil becomes transparent too earlier, leading to failure of both ROM and CSE occurrences.

At the last stage, the compressed nanobunch becomes transparency, leading to heating and decompression of the nanobunch. This makes CSE substantially less efficient at later cycles, resulting in the production of a single (or few) bright CP attosecond pulse.

3. Simulation results

To verify our scheme, 1D PIC simulations are carried out using the code EPOCH [39]. The simulation box is from $-50\lambda$ to $50\lambda$ with a spatial resolution $\lambda/10000$ and 1000 electron and ion macroparticles per cell. The normal incident linearly $P$-polarized and $S$-polarized lasers propagate from the left and right boundary, respectively. Two lasers have the same intensity of $I_p = I_s = 2.1 \times 10^{20} \text{ W cm}^{-2}$, the same wavelength of $\lambda = 800 \text{ nm}$, a Gaussian-temporal profile and a duration of full width at half maximum 18 fs. The density of the fully ionized target is $n_e = 75n_c$ with a thickness of 40 nm. The $S$-polarized laser peak arrives the plasma surface at $(t = 65.50T_0, x = 0)$.

Figure 1(a) plots the electron number density of a thin target versus time irradiated normally by the two counter-propagating linearly $P$-polarized and $S$-polarized lasers where three stages (I–III) as we discussed are shown. The stage (II) is at $65.50T_0$–$66.50T_0$. Because the $S$-polarized laser, which irradiates the foil from the right-hand side, has an initial carrier wave phase delay of $\pi/2$ in our simulations, the isolated attosecond pulse can be obtained in the right direction. The time-frequency spectra of the radiation $E_y$ and $E_z$ at the right side of the target are shown in figures 1(b) and (c), which show that there exist two radiation in a cycle of the drive laser and only one intense radiation is with a wide spectrum. Around the time $t = 65.50T_0$, the transmissivity of the $P$-polarized laser increase fast as shown in figure 1(d), which illustrates that this electron bunch becomes transparency with the emission of an intense radiation driving by both lasers. To obtain an isolated CP attosecond pulse, we choose frequency band 20th–60th corresponding to the green dashed ellipse in figures 1(b) and (c). The waveform of the electric field vector of the CP attosecond pulses is shown in figure 1(e). The symbols ‘A’, ‘C’ and ‘B’, ‘D’ represent attosecond pulse emitted in stage II propagating to the left and right sides, respectively. In the right side, an intense isolated CP attosecond pulse is generated with an intensity of $2.8 \times 10^{19} \text{ W cm}^{-2}$, a duration of $\sim 71\text{as}$, an ellipticity of 0.95 and a phase difference of $-0.53\pi$. 

![Figure 2. Dynamics of the isolated CP attosecond pulse generation in stage II. (a) Longitudinal ponderomotive forces for CLP (red) and CP (black) driving lasers, respectively. (b) Spatiotemporal distribution of longitudinal Lorenz factor $\gamma_x$ of the emitting electrons in the nanobunch. Ion density is shown in the blue background. Green dashed lines represent the retardation paths and the black arrows point the emission nodes. (c) Electron number densities of the emission nodes corresponding to the attosecond pulses ‘A’, ‘B’, ‘C’ and ‘D’ in figure 1(e). (d) Transverse momentum evolution of the emitting electrons.](image)
Longitudinal ponderomotive forces \((J \times B)_z\) are shown in figure 2(a) for the counter-propagating linearly polarized (CLP) and the CP driving lasers, respectively. For comparison, the incident intensity of CP driving laser is \(I_c = 4.2 \times 10^{20} \text{ W cm}^{-2}\) and the other simulation parameters are the same as the CLP cases. The oscillating term of the drive force in our scheme overcomes the limit that CP drive laser cannot produce radiation as the only non-oscillating term. Furthermore, the strong CSE can occur when the density of the electron bunch is at a peak and the bunch moves with a large longitudinal Lorentz factor. Figure 2(b) shows the spatiotemporal distribution of the longitudinal Lorenz factor \(\gamma_x\) of the emitting electrons in the nanobunch which oscillates around the ion background. Green dashed lines represent the retardation paths and the black arrows point the emission nodes. The corresponding electron number densities at the emission nodes for the attosecond pulses ‘A’, ‘B’, ‘C’ and ‘D’ are shown in figure 2(c), respectively. The harmonics order is limited by two factor: (i) the order is proportional to \(\gamma_x^2\) and (ii) the harmonics with a wavelength smaller than the thickness of the electron bunch \((\lambda_0 < d)\) cannot be coherent. For the pulse ‘B’, the emitting electrons in the bunch are with a value of \(\gamma_x^B = 2.1\) and the corresponding electron bunch is with a thickness of \(d_0 = 4\) nm, a density of \(n_{e0}^B = 267 n_c\). While for pulse A, the bunch is with a lower value of \(\gamma_x^A = 1.3\), a thickness of 10 nm and a density of 182\(n_c\). After the pulse B, the target begins to be transparent causing that there is no obvious density spike with a larger value of \(\gamma_x\) at the same time for the pulse ‘C’ and ‘D’. Therefore, at the transparent cycle, the intensities of pulses with a wider and higher filtering band satisfy \(P_{\text{auto}}^B \gg P_{\text{auto}}^C\) and \(P_{\text{auto}}^B \gg P_{\text{auto}}^D\). As shown in figure 2(d) at the radiation point, the transverse momentum reaches its minimum consistent with the theory analysis. Moreover, according the equation (4), the phase difference of the radiation is related to the phase space of the electrons. For the pulse ‘B’, \(p_T\) reaches maximum with \(p_z = 0\), and for the pulse ‘A’, \(p_T\) reaches maximum with \(p_z = 0\). Consequently, an intense isolated CP attosecond pulse is obtained.

To show the robustness of our scheme, a series of 1D PIC simulations are carried out for a wide range of the thickness \(d_0\) and the fluctuation of initial position \(\Delta x\) of the target with a fixed electron number density \(n_e = 75n_c\) as shown in figure 3. The initial phase difference of \(\phi_0 = \pi/2\) and the laser intensity of \(I_p = I_c = 2.1 \times 10^{20} \text{ W cm}^{-2}\) are chosen. The shadow area represents the thickness condition (5) to obtain the CP attosecond pulse in figure 3(a). Relatively stable phase differences approaching \(-\pi/2\) and ellipticity approaching 1 of the attosecond pulse are obtained when the thickness condition is satisfied. When the foil thickness is larger than the condition (5), the ellipticity of the radiation decreases rapidly. Furthermore, the efficiency \(\eta \equiv \int_0^\infty \Delta n |E_x(n)|^2 / \int_0^\infty \Delta n |E_x(n)|^2\) of the obtained attosecond pulse is much higher when the thickness approaches the right side of the condition, which indicates that the target is transparent around the laser peak with an emission of an intense attosecond pulse. Here, \(n\) is the filtering harmonic order, \(n_1 = 20\) and \(n_2 = 60\). Besides, to consider the experimental feasibility of the proposed scheme, the influence of \(\Delta x\) which affects the relative phase of the two ponderomotive forces acting on the combined bunch is shown in figure 3(b). It is clear that CP radiation with high generation efficiency and ellipticity can be obtained in a wide range of \(\Delta x\) except for \(\Delta x = m\lambda/4 + \lambda/8\) where the combined driving force is close to zero. Here, \(m\) is an integer.

To further take the multi-dimensional effects into account, 2D PIC simulations are also carried out as shown in figure 4. Left incident \(P\)-polarized laser and right incident \(S\)-polarized laser irradiate the target normally with the same transverse waist width of \(R = 3\lambda\). The other laser and the target parameters are the same as the above 1D simulation (figure 1). Because 2D simulations for HHG are computationally expensive, the size of simulation box is \(6\lambda \times 12\lambda\) (24000 \times 48000). The electromagnetic fields are recorded by two probe lines at \(x = \pm 2\lambda\). 150 quasiparticles per cell are put in. Figure 4(a) shows the electron number density before the target becomes transparent, where a dense oscillating electron nanobunch is formed at
the peak area of the incident lasers. Besides, due to the Gaussian profile, only the middle area of the target can be transparent which can be improved by using flat-top lasers. An isolated attosecond pulse for the frequency band 20th–60th is shown in figure 4(b). The waveform of attosecond pulse is shown in figure 4(c) along the peak of the attosecond pulse with a phase difference of $-0.52\pi$, an ellipticity of 0.95, a duration of 75 as and an intensity of $\sim 1.2 \times 10^{19}$ W cm$^{-2}$. The divergence angle of 16th–22nd harmonics is about 13 mrad at full-width at half maximum as shown in figure 4(d). Such divergence can ensure enough brilliance for the far-field detection and application [31, 40].

Further, a more available ultra-thin target ($n_e = 400 n_c$, $d_0 = 7.5$ nm) in experiments is considered and the pre-pulse will make the target expand to a wider width and lower density with areal density conserved [34]. So, the case of $n_e = 75 n_c$, $d_0 = 40$ nm can be seen as the expansion result of the ultra-thin target. As shown in figure 4(e), for the case of $n_e = 400 n_c$, an isolated intense attosecond pulse is also achieved with the same filter used, and the Lissajous curves of the filtered $E_y$ and $E_z$ fields on-axis are shown in the insets of figure 4(f). In figure 4(f), the peak intensity of the filtered attosecond pulse for different target thickness $d_0$ is illustrated with the expansion of the high-density target considered. Even when $d_0$ reaches 200 nm, high-ellipticity attosecond pulse can still be obtained with $\epsilon \sim 0.8$ and $\phi \sim -0.5\pi$ as shown in the inset of figure 4(f). A double plasma mirror can be introduced to satisfy the requirements of the contrast ratio of the driving laser [34].

4. Conclusions

In summary, we have shown theoretically and numerically that intense isolated CP attosecond pulses can be generated from ultrathin foil targets irradiated from opposite sides by two relativistic lasers of orthogonal polarizations, where CSE is achieved at the peak laser intensity cycle. With a proper matching condition between lasers and targets, two relativistically oscillating plasma mirrors formed on both sides of the foil are pushed inside by laser radiation pressures, eventually merging together to form a dense electron nanobunch at the peak laser intensity cycle. This nanobunch undergoes the effects of both laser fields with their polarization orthogonal to each other, resulting in CSE of CP HHG, which, afterwards, becomes transparency, making CSE substantially less efficient at later cycles. Therefore, a single intense CP attosecond pulse can be obtained. 2D PIC simulations show that an intense isolated CP attosecond XUV pulse with an intensity of $1.2 \times 10^{19}$ W cm$^{-2}$ and a duration of $\sim 75$ as can be obtained by two lasers with the same intensity of $2.1 \times 10^{20}$ W cm$^{-2}$.
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Data availability statement

The data generated and/or analysed during the current study are not publicly available for legal/ethical reasons but are available from the corresponding author on reasonable request.

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