Giant half-cycle attosecond pulses

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Half-cycle picosecond pulses have been produced from thin photoconductors when applying an electric field across the surface and switching on conduction using a short laser pulse. The transverse current in the wafer plane then emits half-cycle pulses in a normal direction, and pulses of 500 fs duration and \(1 \times 10^5 \text{ V m}^{-1}\) peak electric field have been observed. Here, we show that single half-cycle pulses with a duration of 50 as and up to \(1 \times 10^{15} \text{ V m}^{-1}\) can be produced when irradiating a double foil target with intense few-cycle laser pulses. Focused onto an ultrathin foil, all electrons are blown out, forming a uniform sheet of relativistic electrons. A second layer, placed some distance behind, reflects the drive beam but lets electrons pass straight through. Under oblique incidence, beam reflection provides the transverse current, which emits intense half-cycle pulses. Such a pulse may completely ionize even heavier atoms. With these developments, new types of attosecond pulses in a normal direction, and pulses of 500 fs duration and \(1 \times 10^5 \text{ V m}^{-1}\) peak electric field have been observed. Here, we show that single half-cycle pulses with a duration of 50 as and up to \(1 \times 10^{15} \text{ V m}^{-1}\) can be produced when irradiating a double foil target with intense few-cycle laser pulses. Focused onto an ultrathin foil, all electrons are blown out, forming a uniform sheet of relativistic electrons. A second layer, placed some distance behind, reflects the drive beam but lets electrons pass straight through. Under oblique incidence, beam reflection provides the transverse current, which emits intense half-cycle pulses. Such a pulse may completely ionize even heavier atoms. With these developments, new types of attosecond pulses will become possible.

A major goal of attosecond science1 is to follow electron dynamics on the atomic timescale (24 as), whether in atoms or solids. Different methods for producing single attosecond extreme ultraviolet (XUV) pulses have been described recently, with this end in mind, and some have been tested experimentally2. These methods rely on the generation of high harmonic spectra, which correspond to trains of attosecond pulses in the time domain. Generating these spectra with few-cycle drive pulses and applying high-pass filters may be used to isolate single spikes. For pump–probe experiments it is essential to provide pulses of sufficient intensity, and surface harmonics generated at solid density plasma surfaces may provide high enough photon numbers3 to meet this requirement.

In the present Letter, we follow a very different route, producing half-cycle attosecond pulses from ultrathin relativistic electron sheets (RES). Similar half-cycle pulses, radiated from current sheets in photoconductors at rest, have actually been observed4 and used in ionization experiments5. Another proposal is to produce disk-like, relativistic electron pulses from a synchrotron and to kick them sideways with a magnetic field. Picosecond 200 MW half-cycle pulses are predicted to be produced in this manner6.

The present work is based on two recent developments, the first being the generation of few-cycle high-contrast laser pulses at intensities exceeding \(1 \times 10^{19} \text{ W cm}^{-2}\) (ref. 7). By making use of optical parametric chirped-pulse amplification techniques7 and plasma mirrors8, contrast ratios greater than \(1 \times 10^{10}\) (subpicosecond time intervals, as required in this case) have been achieved. High contrast is crucial so that the ultrathin target foils are not destroyed prematurely. The second innovation concerns the fabrication of these targets, using few-nanometre diamond-like carbon (DLC) foils9, and even mono-atomic carbon layers (graphene)10,11. These foils stay transparent to laser light, even at electron densities of \(n_e \approx 1 \times 10^{24} \text{ cm}^{-3}\) after full ionization. A remarkable feature is that the laser pulse can blow out a major part or even all of the electrons, provided the field amplitude \(E_L\) is larger than the charge separation field \(E_i = e_n a_d / e_0\), where \(e_0\) is the vacuum dielectric constant, \(d\) the thickness and \(e_n a_d\) the charge areal density of the foil blowout. It is the \(\nu \times B_0\) part of the laser interaction that accelerates all electrons in the laser direction. Relativistic energies up to a maximum of \(\gamma = 1 + a_0^2 / 2\) are reached over a few micrometres (where \(a_0 = E_0 / E_0\) is the normalized field strength, \(E_0 = mcω_0 / e\) is the normalizing field, \(ω_0\) is the laser circular frequency, \(e\) and \(m\) are the electron charge and mass, and \(c\) is the velocity of light). The electrons then form a thin RES, separated from the ions and surging on the laser wavefront13.

These high-density sheets may serve as relativistic mirrors, compressing femtosecond probe pulses to attosecond pulses by Doppler shifts of \(4\pi^2\). For this to happen one has to divert the drive pulse from the RES by using an additional foil that is just thick enough (\(\sim 50\) nm) to reflect the light but let the relativistic electrons pass almost unperturbed. Behind the reflector, the electrons propagate in field-free (\(a = 0\)) space, where, due to conservation of canonical momentum, their transverse momentum \(p_t = m^2 c^2 = a\) also vanishes. This means that the RES moves precisely in the normal direction, that is, in the direction of the incident drive pulse. This makes it a perfect relativistic mirror for coherent Thomson backscattering14.

Figure 1 | Scheme of target interaction and half-cycle emission. A three-cycle laser pulse, obliquely incident on a double foil target, is shown at different times, with yellow arrows indicating propagation direction: (i) when starting electron blowout from the production foil, (ii) while diverted by the reflector foil, and (iii) after reflection. The RES (green) passes the reflector in the direction of the green arrows, picking up transverse momentum (small green arrows) due to interaction with reflected light. The transverse current radiates the half-cycle attosecond XUV (HCX, red) pulse propagating in front of the electron layer in the direction of the incident light.

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Here, we consider the same configuration, but now for oblique incidence of s-polarized laser light. A schematic drawing is given in Fig. 1 and a corresponding two-dimensional particle-in-cell (PIC) simulation in Fig. 2. The amazing new feature is the appearance of a strong half-cycle attosecond XUV (HCX) electromagnetic pulse, preceding the RES after emerging from the reflector layer. This is the central new result of this Letter. As we shall explain in the Methods, the HCX pulse results from conservation of canonical momentum under the condition of oblique incidence. Owing to interaction with the reflected laser pulse, each RES electron gets a transverse kick (p_y = pm c tan \theta) in the y-direction of the RES plane. The corresponding transverse current radiates the unipolar half-cycle pulse.

We find HCX formation for a wide range of target and laser parameters, provided they are sufficient to produce the relativistic electron sheet. Here, we present simulations of a reference case to demonstrate the essential features. An s-polarized, three-cycle, \sin^2-shaped laser pulse hits a production foil corresponding to nanomesh graphene. Detailed parameters are given in the Methods. Complete separation of electrons from ions creates an electrostatic field of \( E_0 = 0.76 \text{ TV m}^{-1} \); this is about two orders of magnitude smaller than the laser electric field of \( E_L = 80 \text{ TV m}^{-1} \). Figure 2a shows how this laser pulse drives the electrons out of the production foil (left panel) and how they pass the reflector in the x-direction and form a thin sheet. RES density is plotted for times 7\( \tau_L \), 14\( \tau_L \), and 21\( \tau_L \), where \( \tau_L = 2\pi/\alpha_L \) is the laser cycle. The electron sheet is accompanied by a strong unilateral \( E_y \) pulse, plotted in the right panel of Fig. 2a at times 14\( \tau_L \) and 21\( \tau_L \).

Figure 2b, most importantly, presents line-outs of both electron density \( n_e \) (green) and \( E_y \) field (red) along the central line (y = 0) at times 21\( \tau_L \) and 60\( \tau_L \). It shows how the HCX pulse emerges and gradually separates from the electron layer. The maximum field and width of the HCX pulse are also plotted in Fig. 3a as they evolve along the propagation coordinate. At \( t = 21 \tau_L \), the HCX pulse reaches a peak electric field of \( E_y/E_0 = 1.75 \), corresponding to \( E_L = 7 \text{ TV m}^{-1} \) and a peak intensity of \( I_{HCX} = 1.3 \times 10^{19} \text{ W cm}^{-2} \). This peak then slowly falls due to diffraction while the HCX width is still growing, saturating at \( t = 60 \tau_L \). Plotted as electric field, the full-width at half-maximum (FWHM) width of the saturated pulse is 48\( \alpha_L \) as, corresponding to \( \sim 24 \alpha_L \) as when plotted as intensity. This is the atomic time unit. At \( t = 60 \tau_L \), the relativistic electron sheet has broadened strongly due to Coulomb expansion, and the peak density, initially 30\( n_e \), has fallen to 0.2\( n_e \).

The HCX pulse has now almost separated from the RES and is propagating through vacuum. A remarkable feature is the very sharp front edge, rising from zero to peak values within a few tens of attoseconds. In the present simulation the pulse carries a total energy of 55\( \mu J \), which is \( \sim 1 \times 10^{-4} \) of the incident laser energy. We have
checked that this pulse, \( \sim 7 \) nm long, keeps its half-cycle character over almost a millimetre, before it converts into a single-cycle pulse due to diffraction loss of the long-wavelength components.

Let us now discuss the build-up of transverse current and HCX emission in more detail. The reflector foil plays the key role in switching electron momenta. This is documented in Fig. 3b, which shows \( \gamma, \gamma_y, p_y/mc \) and \( p_y/mc \) as a function of RES position \( x \) along the central line \( y = 0 \). The vertical dashed lines mark the positions of the production and reflector foils. It is seen that, during the initial acceleration phase, RES electrons follow qualitatively relativistic single electron motion, which satisfies \( \gamma - 1 = p_y/mc = (p_y/mc)^2/2 \) and \( p_y/mc = 0 \) when driven by a planar laser pulse polarized in the \( x \) direction.\(^{\text{19}} \) When passing the shaded region close to the reflector foil, \( p_y \) falls to zero, as the laser amplitude \( a_0 \) does, and \( p_y \) pops up, approaching the predicted value of \( p_y/mc = \tan \theta = 0.577 \). This behaviour is derived in the Methods. Inside the dense reflector plasma, the transverse current due to \( p_y \) is screened, but as soon as it emerges from the rear side, it radiates the HCX pulse, while \( p_y \) decreases due to radiation damping. At this point, the present case differs from photoconductive processes, which are mainly emitted during the build-up phase of the transverse current and have opposite polarity. Figure 3b also plots \( \gamma_y = (1 - p_y^2)^{-1/2} \), which is related to the full \( \gamma \) by \( \gamma = \gamma_y[1 + (p_y/mc)^2 + (p_y/mc)^2]^{1/2} \). One might expect that \( \gamma_y \) stays constant behind the reflector because there is no further acceleration by the drive laser. However, there is deceleration due to the charge separation field \( E_z \), which still acts behind the reflector. In particular, fast electrons facing the reflector are decelerated, while front electrons cruise at constant speed. This causes expansion, and a \( \gamma_y \) distribution develops over the sheet. What is actually plotted in Fig. 3b are averages over these distributions.

The electron sheet, moving with \( \gamma_y \) in the \( x \)-direction and having superimposed transverse momentum \( p_y \), emits electromagnetic radiation proportional to the time derivative of the transverse current density in the sheet. In the Methods, the basic equation is solved for the idealized case of a uniform zero-thickness RES. Of course, such treatment does not account for finite-size RES dynamics, but provides useful estimates and scalings. The HCX peak electric field is obtained as \( E_{\text{max}} = \gamma_y a_0 E_0 \sin \theta = 11 \) TV m\(^{-1} \) for the reference case with the separation field \( E_s = 0.76 \) TV m\(^{-1} \) and \( \gamma_0 = 30 \) (Fig. 3b). It should be compared with the simulated value \( E_{\text{max}} = 7 \) TV m\(^{-1} \) at \( t = 21.71 \). In the zero-thickness approximation, the HCX has an infinitely sharp front and then decays exponentially with time constant \( T \approx mc/(\gamma_y e E_0) \approx 75 \) as, corresponding to an FWHM width of 52 as, which is close to the 48 as found in Fig. 2b.

Finer details of the HCX front edge and also HCX growth along the propagation axis (observed in Fig. 3a) are related to the finite size of the RES when emerging from the reflector. In the simulated density profile at this time, the RES shows two peaks at a distance of \( \Delta x \approx 0.008 \Lambda_x \) and some precursor foot extending over \( \Delta x \approx 0.03 \Lambda_x \). The precursor maps directly into the HCX front edge rising over almost the same distance \( \Delta x \). The propagation distance for saturating the peak HCX field can also be estimated as \( L_{\text{sat}} = \Delta x/(1 - \beta_x) \approx 2 \gamma_x^2 \Delta x \approx 14 \Lambda_x \), in fair agreement with Fig. 3a. Here \( L_{\text{sat}} \) is the distance over which the light signal emitted from the second density peak has to travel to catch up with the first peak. These two major HCX contributions then add up coherently and form the peak of the pulse. Owing to retardation, the front contribution is already somewhat damped, and this may explain why the peak field of \( 7 \) TV m\(^{-1} \) is lower than the model result of \( 11 \) TV m\(^{-1} \). It is important to notice that the HCX contributions emitted from different RES layers add up coherently. The coherence is the reason why a major part of the energy deposited in the transverse current is actually radiated, even when the RES has substantially expanded.

We have checked that the results presented in this Letter are robust against changing laser and target parameters. Choosing thicker production foils (DLC with \( n_l = 300 \), and thickness of \( 2-8 \) nm) and Gaussian-shaped laser pulses (FWHM of three cycles) rising in amplitude from \( a_0 = 0.01 \) to peak values of \( a_0 \approx 50 \) and 100, we still find similar HCX pulses (see Supplementary Information). These differences are mainly due to the fact that only a fraction of the electrons of these thicker foils separate from the ions. For the case of a \( 4 \) nm DLC foil and \( a_0 = 100 \), \( \sim 10\% \) of the electrons are found to form the RES. This generates a separation field of \( E_s/E_0 \approx 1 \), to be compared with a 100 times larger driving field as in the reference case of this Letter. An HCX pulse is produced with a peak field of \( E_s/E_0 \approx 2 \) and width of 169 as.

In conclusion, we have found an efficient new option to generate single attosecond pulses. This makes use of a double foil target to produce and purify a RES. Under oblique incidence, it radiates a half-cycle pulse with a duration of a few \( 10 \) as, a peak electric field up to \( 1 \times 10^{13} \) V m\(^{-1} \), and pulse energies up to the 0.1 mJ level. The laser-to-HCX conversion efficiency amounts to a few \( 10^{-4} \). Most important for experiments is the extremely sharp front edge with rise times of a few tens of attoseconds. HCX properties and scaling relations have been derived in simple analytical terms. Compared to Thomson scattering as a way to generate single attosecond pulses\(^{18} \), the present method is simpler to implement because only a single laser beam is required. It is hoped that the present Letter will stimulate experiments.

**Methods**

**Transverse momentum.** For oblique incidence, electron dynamics close to the reflector are quite complex. A simple analytic approach is possible when transforming to a frame in which the laser beam is vertically incident, as depicted in Fig. 4 (ref. 16). By first rotating the \( x'y'z' \) layer frame to the \( x'y'z \) layer frame and then applying a Lorentz boost in the \( y \) direction with \( \beta_y = -\sin \theta, \gamma_y = \sec \theta \), electron energy-momentum \( (\gamma_p, p_x, p_y, p_z) \) transforms to \( (\gamma_p, p_x, p_y', p_z') \) with \( \gamma' = (1 + \beta_y^2 + p_y'^2 + p_z'^2)^{1/2} \); here energy is in units of \( mc^2 \) and momenta in \( mc \). Owing to the conservation of canonical momentum, the \( z \)-component of momentum (vertical to the plane of Fig. 4) is invariant and equals the vector potential \( a_z \) for s-polarized light. \( p_z = p_z'' = a_z \). In the boosted frame, the \( y \)-component \( p_y'' = \beta_y \gamma_y \) is also conserved. Transforming back to the laboratory frame, we find

\[
p_y'' = p_y' + (p_y'' - \gamma') \sin \theta
\]

where \( p_y'' = \gamma' = (1 + \tan^2 \theta + p_z''^2 + a_z^2)^{1/2} \). Here we notice that, behind the reflector where \( a_0 = 0 \), we have \( p_y'' = \gamma' = 0 \) for relativistic electrons, provided that \( \beta_y \cos \theta \gg 1 \). In this case, \( p_y'' = \tan \theta \) stating that electrons emerge from the reflector with non-zero momentum in the direction of the RES plane. Momentum transfer into the \( y \)-direction occurs within a narrow region close to the reflector (Fig. 3b), where \( a_0 \to 0 \).

**Scaling of HCX emission.** HCX generation from a RES can be described by the one-dimensional wave equation

\[
(\partial^2/\partial x^2 - c^2 \partial^2/\partial t^2)E_x = e_x^2 c^2 J_{\text{sat}}/\partial t
\]

where \( J_{\text{sat}} = -e_0 j_{\text{sat}} \) is the radiating current. The solution can be written as

\[
E_x(t, x) = -(2a_0)^{-1} \left[ \text{J} \left( t - \left( x - \frac{1}{2}c t \right) \right) \right] e^{i \Delta x / \Delta t}
\]

where \( \text{J}(t) \) is a complicated function of space and time in general. Here, we restrict ourselves to an RES of zero thickness, approximating the density profile as \( n(x) = n_0(\delta(x/d)) \) with finite area.
density \( \rho(x) \) \( dx = \rho_0 \), and first consider the rest frame (index \( R \)), choosing uniform velocity \( \mathbf{v}_0 = 0 \) and \( \beta_0 = \sin \theta \) initially. The solution consists of two half-cycle electromagnetic pulses emitted symmetrically in the \( x^{-} \)-direction, just as observed in the photocathode-conductor experiments. Expressing \( I_0 \) in terms of transverse momentum \( p_y \) for the wave travelling in the \( x^{+} \)-direction

\[
E_y(t_y) = (E_x/2)p_y/(\gamma t_y)/mc
\]

where \( t_y = t-x/c \) and \( E_x = cE_0(x_0) \). The relativistic factor \( \gamma = (1 + (p_y/mc)^2)^{1/2} \) can be estimated as \( \gamma \approx 1 \). At \( x = 0 \) (that is, \( t_y = t \)), where the layer lies, \( E_x \) damps the electron momentum according to \( dp/dt = -eE_x(t) \). One finds that both \( p_y(t) \) and the emitted pulse \( E_y(t) \) decay exponentially on a timescale \( \tau_E = 2mc/(eE_x) \). These results apply to the rest frame of the layer. Performing a Lorentz transformation to the laboratory frame, in which the layer is moving with velocity \( \beta \), and \( \gamma = (1 - \beta^2)^{-1/2} \) in the normal direction, we find the coordinate of the forward HCX pulse \( \tau = \gamma (1 - \beta) t_0 = t_0/(2\gamma) \) and the electric field \( E_x = \gamma (1 + \beta)^2 E_y \). The maximum HCX field scales as \( E_x \approx \gamma E_y \sin \theta \) and the pulse width as \( T \approx (\gamma^2 E_y) \).

**PIC simulation.** We carried out all simulations using the JPIC code\(^17\), which applies a field solver free of numerical dispersion in the \( x \)-direction\(^19\). Because an isolated HCX pulse contains extremely broad frequency components, for accurately used them in our PIC simulation.

(\( \gamma(x) \approx \gamma \sin \theta \))

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**Author contributions**

H.-C.W. discovered the new effect described in this Letter, carried out all simulations and developed the basic theory. J.M.-t.-V. wrote the paper and clarified some details of the physics. Both authors take full responsibility for the presented results.

**Additional information**

The authors declare no competing financial interests. Supplementary information accompanies this paper at www.nature.com/naturephotons. Reprints and permission information is available online at http://www.nature.com/reprints. Correspondence and requests for materials should be addressed to H.-C.W.