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Measurement of the phase of few-cycle laser pulses

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Abstract. Attosecond science can take enormous advantage of intense laser pulses consisting of two optical cycles or less (few-cycle pulses). The temporal variation of the field, which directly governs strong-field interactions, therefore depends on whether the maximum of the pulse amplitude coincides with that of the wave cycle or not, i.e. on the phase of the field with respect to the pulse envelope. It is demonstrated that the direction of electron emission from photoionized atoms can be controlled by varying the phase of the field, providing a tool for its accurate determination. Different phase-dependent effects in strong-field photoionization and their suitability for phase measurement are discussed as well as the limitations of the *f*-to-2*f* scheme.

1. Attosecond science and the 'absolute' phase

Investigation of phenomena elapsing on the attosecond scale take advantage of an intriguing effect discovered many years ago in strong-field laser atom interaction: electrons produced by photoionization may be driven back to the vicinity of the ion core—typically within less than an optical cycle—and interact there again [1, 2]. Recombination leads to generation of coherent soft X-ray radiation [3] and eventually attosecond pulse generation [4, 5]. Elastic rescattering boosts the electron energy [6] and inelastic scattering is at the basis of correlated double ionization [7].

The trajectories of the revisiting electrons are precisely locked to the laser field and are launched in time intervals confined to the sub-femtosecond domain. However, this process repeats itself in every optical cycle (or even half-cycle). This and—even worse—multiple returns of an electron wave packet to the ion core question the suitability of the phenomenon as a general approach to attosecond physics. The most convincing answer to the problem obviously is to confine the driving laser field to more or less a single optical cycle (few-cycle pulses)—see figure 1 for example. In fact, this was achieved years before attosecond pulses were

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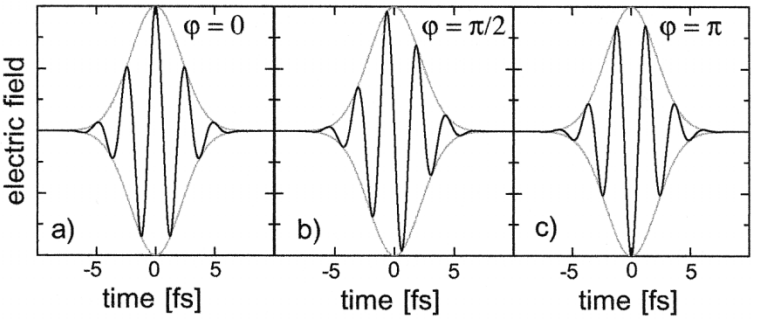

Figure 1. 5 fs laser pulses consist of less than two optical cycles (few-cycle pulses). The temporal evolution of the electric field of such pulses depends on the phase of the carrier with respect to the envelope, the so-called carrier-envelope (CE) phase or 'absolute' phase. By convention, the origin of the time scale is chosen to be at the maximum of the envelope of the pulses. Then an 'absolute' phase $\varphi = 0$ corresponds to cosine-like pulses and $\varphi = -\pi/2$ to sine-like pulses.

measured for the first time [8]. Few-cycle pulses, however, suffer from a novel problem: the temporal variation of the laser field and thus attosecond pulse generation depend on the 'absolute' phase. It is exactly this problem which will be addressed in this review and it will be shown that the very electron trajectories inducing attosecond pulses can also be used to measure the 'absolute' phase.

Generally, the electric field of a pulse can be written as

$$\mathcal{E}(t) = \mathcal{E}_0(t)\cos(\omega t + \varphi). \tag{1}$$

$\mathcal{E}_0(t)$ denotes the envelope of the pulse which—for few-cycle pulses—varies barely slower than the laser frequency $\omega/2\pi$. For a definition of the absolute phase φ we use the convention that the pulse has its maximum at $t = 0$. Accordingly, pulses with $\varphi = 0$, $-\pi/2$, π and $\pi/2$ are dubbed cosine-like, sine-like, $-$cosine-like and $-$sine-like. In order to fix also the spatial alignment of the pulse, horizontally polarized radiation is assumed in this paper. Further, the convention is assumed that, looking in direction of pulse propagation, the global maximum of a cosine-like pulse points to the right.

Absolute-phase effects were detected [9] even before it became possible to stabilize the absolute phase. The latter came from an ingenious invention in frequency metrology [10–13] and now allows one to collect data with thousands of pulses of identical shape [14]. Shortly after that, the phase was measured and the spatio-temporal shape of the pulse unambiguously determined [15, 16].

2. Experimental set-up and principles of operation

Our approach for phase measurement is based on the conjecture that the spatial symmetry of photoelectron emission from a gaseous (i.e. isotropic) medium should be governed by the symmetry of the pulse. Since the average field of any propagating pulse is zero, a nonlinear dependence of photoelectron yield on intensity is required. It is however *not* essential to assume an extremely strong nonlinearity or even a threshold-like dependence. Rather, the interplay of

nonlinear photoionization and deflection of the photoelectrons in the asymmetric field leads to an asymmetric yield [17].

Based on this reasoning an instrument consisting of two opposing field-free time-of-flight (TOF) photoelectron spectrometers was devised. Due to its characteristic appearance it was eventually called a stereo-TOF spectrometer, see figure 2. The electron detectors can accept several electrons per laser pulse. The TOF is measured by computer-hosted multiscalers. Operation of the instrument in event and histogram mode is possible. The former allows seeking effects of the absolute phase even if it changes in a completely random way from laser pulse to laser pulse. This is still possible under situations where only a few electrons are detected per laser pulse, i.e. too few electrons to see phase effects from a single laser pulse: the phase-induced asymmetry of photoelectron emission translates into an anticorrelation of electrons emitted to the left and to the right [9].

Phase-stabilized pulses allow one to operate the instrument in histogram mode and to record photoelectron spectra. The experimental set-up offers two possibilities to change the absolute phase in a controlled way. Both rely on the fact that the envelope propagates with the group and the carrier wave with the phase velocity. Differences of both velocities thus will induce a change of the absolute phase.

The first possibility exploits the fact that group and phase velocity are different for any focused wave. This is known as the Gouy effect [18]. For

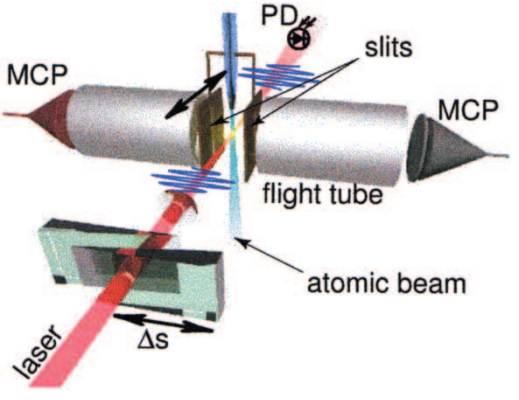

Figure 2. 'Stereo-ATI' spectrometer. Two opposing electrically and magnetically shielded time-of-flight spectrometers are mounted in an ultrahigh vacuum apparatus. Rare gas atoms fed in through a nozzle from the top are ionized in the focus of a few-cycle laser beam. The focal length is 250 mm (the lens shown in the sketch is in reality a concave mirror) and the pulse energy was attenuated to 20 μJ. The laser polarization is linear and parallel to the flight tubes. Note that the laser field changes sign while propagating through the focus. Slits with a width of 250 μm are used to discriminate electrons created outside the laser focus region chosen. The slits can be moved from outside the vacuum system. A photodiode (PD) and micro-channel plates (MCP) detect the laser pulses and photoelectrons, respectively. The laser repetition rate is 1 kHz and 50 electrons per pulse are recorded at each MCP. A pair of glass wedges (apex angle 2.8°) is used to optimize dispersion and adjust the absolute phase. Therefore, movement of one wedge by Δs introduces $\Delta x = \Delta s \tan(2.8°)$ glass.

a continuous-wave Gaussian beam the phase changes along the axis of propagation z as

$$\varphi = \arctan\left(\frac{z}{z_R}\right), \tag{2}$$

where z_R is the Rayleigh length of the laser focus. Non-trivial modifications of equation (2) are necessary for few-cycle pulses [19]. In our experiment a pair of opposing grounded plates with narrow slits transmit only electrons originating from a region of the focus much smaller than z_R. Moving this pair of slits with respect to the focus leads to photoelectrons created by pulses with different absolute phase.

The second possibility is based on glass dispersion. A piece of glass of thickness d induces a change $\Delta\varphi$ of the absolute phase given by

$$\Delta\varphi = \omega d\left(\frac{1}{v_{ph}} - \frac{1}{v_g}\right), \tag{3}$$

where v_{ph} and v_g are phase and group velocity, respectively, given by

$$v_{ph} = \frac{c}{n}, \tag{4}$$

$$v_g = \frac{c}{n} + \frac{c\lambda}{n^2}\frac{\partial n}{\partial \lambda}. \tag{5}$$

This means that, at the central wavelength $\lambda_0 = 760\,\text{nm}$ of the laser used in our experiments, a fused silica glass plate with a thickness $d = 26\mu\text{m}$ changes the absolute phase by $\Delta\varphi = \pi$. A cosine-like pulse would be converted into a minus-cosine-like pulse.

3. Comparison with the *f*-to-2*f* scheme

Currently, stabilization of the absolute phase of amplified few-cycle pulses is achieved with a two-stage servo-loop using two *f*-to-2*f* phase detectors following the oscillator and the amplifier. While strong-field photoionization is directly sensitive to the absolute phase, systematic errors have to be expected from phase measurements with a *f*-to-2*f* set-up, because it relies on comparing the phases of the blue and red wing of the spectrum rather than on its centre of mass. In order to explain this, we assume that an amplifier is seeded by a phase-stabilized oscillator only, i.e. there is no stabilization for any possible drifts in the laser system after the oscillator. The absolute phase after the amplifier indeed drifts as expected. For simplicity, a change in phase of $\Delta\varphi = 2\pi$ is assumed to occur in one minute. This is compatible with experimental observations. Here, we do not specify the origin of the phase drift. Now we add phase stabilization after the amplifier and compare two phase stabilization schemes with phase detectors using the *f*-to-2*f* method and direct measurement of the absolute phase as proposed below. Either of these schemes would act on a glass wedge and thus change the amount of glass in the laser beam by d_{1f-2f} and d_{abs} which would induce phase changes such that $\Delta\varphi$ is compensated.

For the *f*-to-2*f* method, $\Delta\varphi$ can be calculated using the travel times through the glass plate $\tau_{red} = d_{1f-2f}\, n_{red}/c$ and $\tau_{blue} = d_{1f-2f}\, n_{blue}/c$ for the red and the blue wing of the spectrum generated through white light generation.

$$\Delta\varphi = 2\pi \frac{\tau_{red} - \tau_{blue}}{T_{blue}}. \tag{6}$$

T_{blue} is the optical period of the modes in the blue wing which have to be used here because the red modes are frequency doubled. This leads to

$$\Delta\varphi = 2\pi \frac{d_{1f-2f}}{\lambda_{blue}}(n_{red} - n_{blue}) \quad \text{or} \quad d_{1f-2f} = \frac{\Delta\varphi}{2\pi} \frac{\lambda_{blue}}{n_{red} - n_{blue}}. \tag{7}$$

For techniques directly sensitive to the absolute phase, the respective expression can be derived from equation (3) and by using the first-order Taylor expansion of $1/v_g$:

$$\Delta\varphi = 2\pi d_{abs} n'(\lambda_0) \quad \text{or} \quad d_{abs} = \frac{\Delta\varphi}{2\pi n'(\lambda_0)}\,, \tag{8}$$

where $n' = \partial n/\partial\lambda$. In general, d_{1f-2f} and d_{abs} are different, and the difference depends on the central wavelength λ_0 of the few-cycle pulses and the spectrum generated through white light generation. In particular the latter can be difficult to control experimentally. While λ_{blue} typically is in the range of 550 nm, deviations of tens of nanometres are possible. Figure 3 displays the difference of d_{abs} and d_{1f-2f} as a function of λ_{blue} and for some λ_0 for a phase difference $\Delta\varphi = 2\pi$ to be corrected.

This analysis reveals that phase stabilization schemes using these two phase detectors would introduce different amounts of glass in response to a phase change

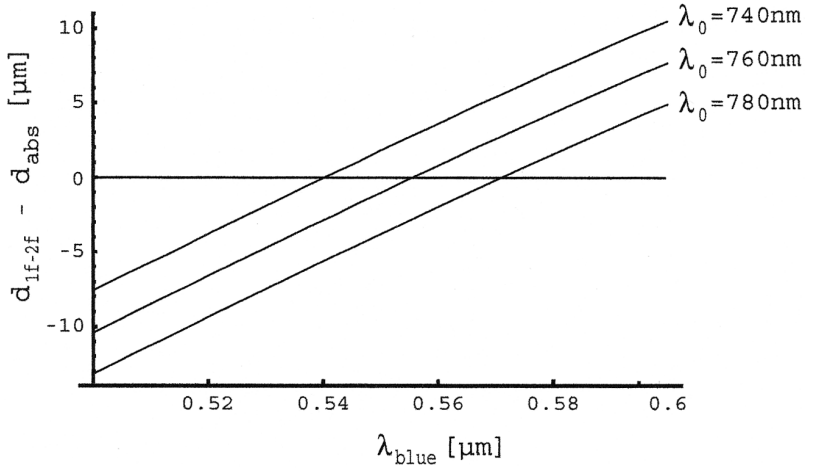

Figure 3. Two phase stabilization servo loops using the *f*-to-2*f* method or a direct measurement of the absolute phase are supposed to correct a phase drift eventually reaching 2π by introducing fused silica glass of thickness d_{1f-2f} and d_{abs}, respectively, into the laser beam. Depending on the centre wavelength λ_0 of the spectrum of the few-cycle laser and the frequency 2*f* (corresponding to λ_{blue}) used by the *f*-to-2*f* interferometer, both servo loops will introduce a different amount of glass into the beam. It should be noted that the dispersion of 26 µm fused silica changes the absolute phase by π.

$\Delta\varphi$. Even for a modest phase change of $\Delta\varphi = 2\pi$ this difference can easily be several micrometres of glass, implying a systematic error of several percent for the *f*-to-2*f* scheme. The weak point of this analysis certainly is that the origin of $\Delta\varphi$ was not specified and the same phase change is assumed for the direct and the *f*-to-2*f* phase measurement. For example: in the trivial case where $\Delta\varphi$ would be caused by the same glass as the one by which $\Delta\varphi$ is to be compensated later, there would be no difference between both phase detectors. This, however, is a special case of no relevance to the problem considered. For an estimate of the magnitude of the systematic deviations to be expected from the *f*-to-2*f* scheme, the above assumption is reasonable.

4. Phase-dependent photoelectron spectra

We now turn to the experiment measuring the absolute phase directly. Figure 4 shows a series of photoelectron spectra recorded for different positions of the fused

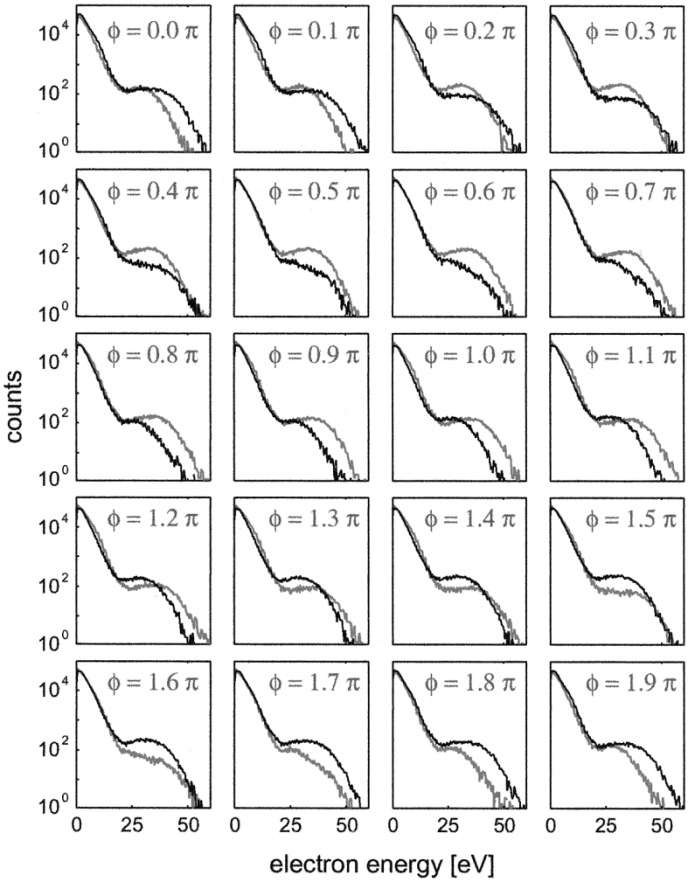

Figure 4. ATI spectra of xenon (semi-logarithmic scale) measured with the left (grey curves) and right (black curves) detectors with phase stabilization. The spectra were acquired by integrating 40 s at each position and by changing the phase in constant steps of $\pi/10$. The labels indicate the absolute phases corresponding to each spectra, as deduced from the phase assignment discussed below.

silica wedge shown in figure 2. While the value of the absolute phase is not known *a priori*, the phase difference between the spectra can be easily calculated using equation (3).

The spectra immediately reveal a remarkable dependence on the absolute phase. Obvious checks are that a phase difference of π should swap left and right and that a phase difference of 2π restores the original spectra. (Here we disregard small effects caused by changes of the pulse duration due to changes in second-order dispersion.) Other observations are as follows.

(1) The phase dependence of the high-energy electrons is much stronger than the one for low-energy electrons. However, it should be emphasized that figure 4 displays the spectra on a semi-logarithmic scale. In fact, figure 5 reveals that the low-energy electrons are also phase dependent.
(2) There is *no* absolute phase for which the spectra emitted to opposite directions would be identical. At first glance, this result might appear surprising. More importantly, however, it means that a single pair of spectra can provide an unambiguous measurement of the absolute phase.

Closer inspection of the spectra reveals that there are at least four phase-dependent effects each of which can be used for an independent measurement of the absolute phase. Before discussing these effects in more detail, a few remarks about the nature of photoionization with intense ultrashort laser pulses are necessary.

At high intensities atoms may absorb more photons than necessary for ionization. Therefore, electron energies exceeding the photon energy by far can be observed [20]. The effect is known under the (misleading) name above-threshold ionization (ATI). Atoms with high ionization potential (i.e. rare gas atoms) can be exposed to intensities exceeding $5 \times 10^{13}\,\mathrm{W\,cm^{-2}}$. Then the spectra develop a plateau-like annex at high energy [6]. Such spectra are striking evidence for non-perturbative behaviour. The basic mechanism was identified shortly after the discovery of the ATI plateau: electrons may be driven back to the vicinity of the parent core and absorb additional photons upon rescattering [21, 22].

5. Four ways for phase measurement

5.1. *The total electron yield*

The simplest phase effect is that of the total electron yield, i.e. the spectra as shown above integrated over energy, see figure 5. Although the contrast of this phase effect is rather low (10% for a FWHM pulse duration between 6 and 7 fs), the huge number of electrons provides low noise. An obvious disadvantage is that the phase has to be scanned in order to measure it. Evidently, measuring the left–right ratio for a fixed phase cannot determine the sign of the phase for those phases leading to identical *total* yield in opposite directions. In addition, the left–right ratio of the total yield depends critically on the pulse duration and perhaps on other quantities. Therefore for all other phases too, the left–right ratio of the total yield measured at one fixed phase cannot be used for a phase measurement.

In any case, a sufficient understanding of the physics of strong-field ionization is required in order to assign the absolute phase. Notwithstanding that it has been claimed for ten years that ATI is understood and further research obsolete,

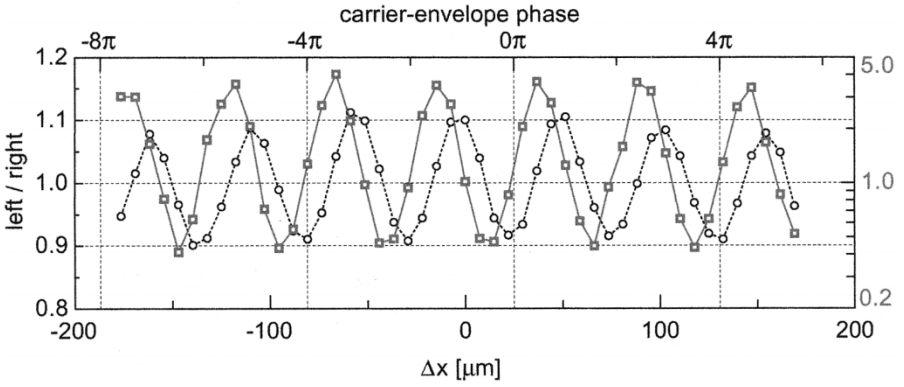

Figure 5. Left–right ratio of the total electron yield (circles) and high-energy electrons (squares) as a function of glass thickness Δx added or subtracted by moving one of the wedges. $\Delta x = 0$ corresponds to optimal dispersion compensation, i.e. the shortest pulses. Maximal left/right ratio for the total yield does not coincide with that for high-energy electrons. Note the different scales for low- and high-energy electrons. The upper x-scale indicates the carrier-envelope phase of the pulse, as deduced from the comparison with theory.

phase-dependent ATI spectra reveal immediately that even the most basic phenomena of strong-field ionization are far from a sufficient understanding.

The only quantitatively reliable method to calculate the phase dependence of strong-field ionization is numerical integration of the Schrödinger equation. This is easily possible for atomic hydrogen only. Presently there is no *ab initio* calculation available for rare gas atoms beyond He. Experimentally, however, the heavy rare gases are much more favourable.

The simplest theoretical approach is the so-called strong field approximation (SFA). It assumes that the electron enters the continuum at some instant t_0 with zero velocity. Depending on t_0, the electron will be accelerated by the oscillating electric field of the laser. The drift energy of the photoelectrons can be easily calculated exploiting conservation of canonical momentum *if* the atomic potential can be neglected:

$$p_{\text{drift}} = -eA(t_0). \tag{9}$$

Here $A(t)$ is the vector potential and e the elementary charge. The dependence of the drift momentum on the vector potential rather than on the electric field implies a counter-intuitive result: the SFA predicts equal yields to opposite directions for cosine-like and maximum asymmetry for sine-like pulses! Considering the high field strength and the correspondingly big oscillation amplitudes of the electrons, the key approximation of the SFA of omitting the atomic potential appears to be reasonable. In fact, the SFA in its classical and quantum version have been able to explain a lot of effects in strong-field laser matter interaction qualitatively. For a measurement of the absolute phase, a quantitative understanding is required. In fact, it turns out that a phase assignment based on the SFA would not not be accurate. This has been confirmed by numerically solving the TDSE [23]. However, the weak asymmetry exhibited by the ratio of the total number of electrons emitted to opposite directions can be explained in the framework of the SFA [17].

5.2. *The plateau electrons*

Obviously, the problems in understanding the phase dependence of the total ionization yield are related to the influence of the atomic potential which is neglected in the framework of the SFA [23]. The total yield is dominated by photoelectrons with an energy close to zero. Clearly, these electrons are affected most strongly by the atomic potential. It is evident that the high-energy electrons in the plateau region of the spectra are more favourable for phase determination. Indeed, it is known that the dominant overall features in this spectral region do not depend on specific properties of the atomic species or on intensity. Another good argument for using the plateau electrons certainly is their higher sensitivity to the absolute phase. The disadvantage however is that more than 100 times fewer high-energy than low-energy electrons are created.

The kinetic energy spectrum of the plateau electrons is sensitive to the absolute phase for several reasons. First of all, high-energy electrons returning to the ion core can only be created in sub-femtosecond time intervals close to peaks of the electric field of the laser pulse. However, the probability that they tunnel through the atomic potential at t_0 depends exponentially on the field strength $\mathrm{E}(t_0)$ and—as few-cycle pulses are involved—it is likely only for those very few optical half-cycles close to the pulse maximum. Generally, the highest kinetic energies are reached for electrons returning to the core at times t_1 when the electric field becomes nearly zero ($\mathrm{E}(t_1) = 0$). For few-cycle pulses in addition, the field amplitude E_0 needs to be as high as possible for $t > t_1$ in order to allow efficient acceleration after rescattering. Since the start time t_0 and return time t_1 differ by almost one optical cycle, both conditions—namely the highest possible field strength at t_0 and highest possible amplitude after t_1—are hard to meet and result in a strong dependence of photoionization on the absolute phase. Number, strength and timing of the wave-packets lead to distinctive structures in the ATI spectra. Their analysis therefore provides detailed information about the key processes of attosecond science. Quantum mechanical calculations using the SFA are in very good qualitative agreement with this classical treatment [24].

Calculating ATI spectra of rescattered electrons along the lines of this model is straightforward: each instant t_0 during the pulse leads to a different electron trajectory, the weight of which is determined by the static tunnelling probability at t_0. The trajectories are given by the equation of motion $\ddot{x} = (e/m)\mathcal{E}(t)$ and the initial conditions $\dot{x}(t_0) = x(t_0) = 0$. Only trajectories returning to the core at some time $t_1 > t_0$ are considered. Upon return elastic rescattering is assumed, i.e. the velocity changes its sign: $\dot{x}(t_1 - \epsilon) = -\dot{x}(t_1 + \epsilon)$. This leads to further acceleration of the electrons in the laser field. Refinements of this model may include wave-packet dispersion and Coulomb refocusing. It should be emphasized again that the energy of these electrons is big as compared to the ionization potential. Therefore, the potential of the ion core will not deflect the electrons appreciably. It is unlikely that electrons predicted to fly to the left for the potential neglected will in reality fly to the opposite direction. This simple model can be extended to include interference effects by calculating the classical action for the trajectories and adding trajectories with the same final momentum coherently.

5.3. *The $2U_p$ cut-off*

For electronic trajectories without rescattering the SFA predicts a maximum drift energy given by $2U_\mathrm{p}$, U_p being the laser's ponderomotive potential. In the

case of few-cycle laser pulses, however, the situation is markedly different. Conservation of canonical momentum (equation (9)) implies that the maximum drift energy of the direct electrons (i.e. the low-energy electrons which did not rescatter) in, e.g. the positive (right) direction is expected for the most negative value of $A(t_0)$. This occurs for a vector potential proportional to $-\cos(\omega t)$. Since $E(t) = -\partial A(t)/\partial t$, this corresponds to an electric-field waveform proportional to $-\sin(\omega t)$, i.e. to an absolute phase $\varphi = \pi/2$.

Experimentally, the direct electrons' cut-off can be observed at intensities higher than $10^{14}\,\mathrm{W\,cm^{-2}}$ by the clear kink appearing in the ATI spectra at electron energies around $2U_p$. This is shown in figure 6 for three different absolute phases. The value of the phase can be determined by the evident asymmetry around $2U_p$. Note that this approach is in excellent agreement with the phase assignment made by means of the asymmetries in the ATI plateau. In fact, the two methods are complementary: sine-like pulses ($\varphi = -\pi/2$ and $\varphi = \pi/2$) are more easily distinguishable by observing the $2U_p$ cut-off, while cosine-like pulses ($\varphi = 0$ and $\varphi = \pi$) are more easily distinguishable by observing the plateau.

5.4. *Phase-dependent interferences*

The typical signature of ATI used to be the regular series of peaks spaced by the photon energy. This peak structure reflects the periodicity of the electromagnetic field. Within the strong-field approximation one would argue that in each optical cycle a wave packet tunnels through the Coulomb potential due to optical field ionization. In the energy domain this leads to the observed ATI peaks.

For few-cycle pulses, however, there might be only a single cycle that produces a wave packet of significant amplitude. Whether one or two cycles contribute to the electron yield at a given electron energy depends on the absolute phase and on the emission direction. Obviously, it is the availability of which-way information which is at the basis of the visibility of peak structure. The situation is analogous to a double-slit arrangement. It is evident that the peak structure offers another

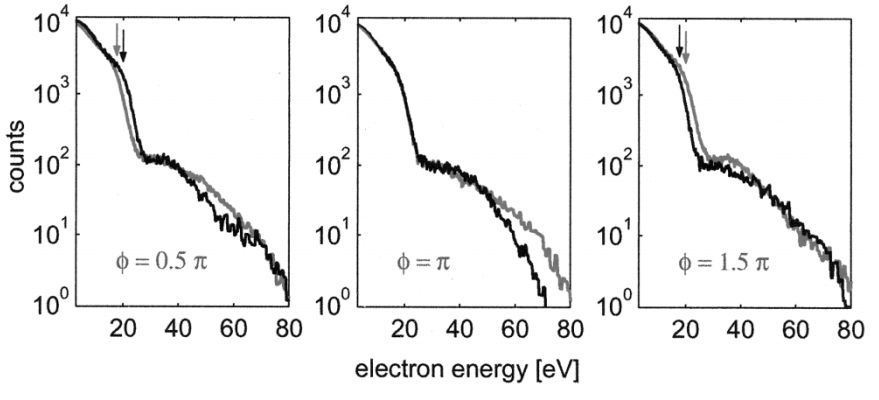

Figure 6. ATI spectra (semi-logarithmic scale) measured in the left (grey lines) and right (black lines) direction at relatively high intensity ($1.5 \times 10^{14}\,\mathrm{W\,cm^{-2}}$). The clear cut-off difference of the direct electrons allows a determination of the absolute phase. Note that, due to the relation between electric field and vector potential, pulses with absolute phases $\varphi = 0$ and $\varphi = \pi$ cannot be distinguished by the direct electron cut-off only. This ambiguity is easily removed by taking into account the high-energy part of the spectrum, too.

independent path for phase measurement. Again, however, a sufficient theoretical understanding is necessary. Preliminary results obtained by the semi-classical model described in section 5.2 indicate a phase assignment compatible with the results obtained by the phase effects described before.

6. Conclusion

We have shown that stereo time-of-flight spectroscopy of above-threshold ionization photoelectrons is the most suitable approach for measuring the absolute phase of few-cycle laser pulses at present. Several phase-dependent effects of the photoelectron spectra have been discussed. The theoretical understanding for the high-energy photoelectrons is sufficiently reliable to allow an unambiguous phase assignment. Consistent results have been found for different phase effects exhibited by the photoelectron spectra.

While this paper has concentrated on the *measurement* of the absolute phase, we would like to emphasize that the significance of phase-dependent ATI is by no means limited to this technical issue. Instead of using the phase dependence of strong-field ionization for measuring the absolute phase, the absolute phase can also be used to control photoionization. In the simplest case, the emission direction of the photoelectrons is controlled. In fact our experiments can be interpreted exactly in this way. It appears to be justified to regard this as a new type of coherent control. More sophisticated experiments in the future could consider dissociation of molecules, in particular if the polarization could be controlled on the few-cycle time scale. This would allow one to deform molecular potentials in a virtually arbitrary way.

It has already been pointed out that the present understanding of strong-field laser atom (and molecule) interaction is deficient. We already know that the phase dependence of the low-energy electrons deviates from theoretical predictions. In contrast to most features observed in strong-field laser atom interaction, this phase-dependence can be measured extremely precisely. Such data, virtually free of detrimental side-effects, such as focal averaging, can provide the experimental evidence on which theoretical modelling beyond qualitative estimates can be founded.

Finally, it should be noted that the ATI plateau electrons are created by the very process underlying all current approaches to attosecond physics. Their analysis, in particular the interference effects mentioned in section 5.4, provides detailed information about the dynamics of the underlying mechanisms.

In light of this, the statement at the beginning of this contribution saying that few-cycle pulses would 'suffer' from a novel problem needs to be revised: it rather appears that the absolute phase as a new parameter of laser pulses will literally open up a new dimension for quantum optics experiments.

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