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Two-Photon Resonance Generation of Single and Double Attosecond Pulses in Ar^+

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Abstract—The resonance generation and propagation of single as well as double attosecond pulses in Ar^+ are presented via numerical simulations. Two-photon Raman-type self-induced transparency (RSIT) of the pulses is investigated. The transient pulse behavior when the input pulse is not in the RSIT regime is investigated to optimize the transition from an input few-cycle pulse to an output attosecond pulse that is in the RSIT regime. The spectral dynamics of the pulses is also presented. The coherent lengths of plasma mismatching are calculated for the attosecond pulses.

1. INTRODUCTION

Attosecond (as) pulses open up many new opportunities for science and technology (see, e.g., [1–5] and references cited therein). For example, they will permit the study of electron dynamics in the attosecond regime, of autoionization states, and of inner-shell electron relaxation processes in medium- and high-Z atoms as well as the attosecond dynamics of nuclear processes. Much attention is being given to the generation and measurement of attosecond pulses using two main methods: first, high-harmonic generation (HHG) in rare gases; and second, high-order stimulated Raman scattering (SRS) at vibrational and rotation transitions in molecules, including solid hydrogen. The advantage of the SRS method is the high-energy conversion efficiency, and the disadvantages are low ionization thresholds in comparison with rare gases as well as problems of propagation. For example, the spectrum of a 300-as unipolar pulse covers a spectrum of more than 10^5 cm^{-1} (~13 eV). Such a pulse (and shorter pulses all the more) will capture very many resonant transitions in, for example, H_2, that have the largest vibrational transition; the fundamental vibrational transition frequency is about 4401 cm^{-1} (the anharmonicity constant is 121.34 cm^{-1}) [6]. As for HHG, the low-energy conversion efficiency ($\leq 10^{-3}$ at XUV and 10^{-5} at X-ray wavelengths) is a problem, even taking into account methods (see, for example, [4, 7–9]) that were proposed recently to increase the efficiency of HHG. Nevertheless, the observation of a train of 250-as pulses that were spaced by 1.35 fs (half the cycle time of the driving, 40 fs; wavelength, 800 nm) from HHG in argon was reported by Paul *et al*. [10] and Agostini and Paul [11]. All of the above-mentioned methods produce discrete spectra and, therefore, only trains of as pulses, whereas it is the generation of single attosecond pulses with high intensities that is of prime interest for different applications, including XUV pump-probe experiments on an attosecond time scale. Thus, substantial effort is being dedicated to generating single attosecond pulses [3–5, 12–14].

A new way, based on electronic stimulated Raman self-scattering (SRSS) of fs pulses in rare gases, to generate intense individual attosecond pulses for the first time was proposed by us [12]. Additionally, some possibilities for generating single subfemtosecond pulses from initial 4.63-fs, 7-fs, and 10-fs pulses using SRSS in singly ionized rare gases was considered in our previous papers [1, 13]. SRSS is resonant coherent SRS when the frequency of the Raman transition is captured by the spectrum of initial pulse, and it was first investigated at vibrational and rotation transitions [15–17]. A continuous Stokes frequency shift with spectral broadening of femtosecond pulses under the conditions of SRSS in vibrational transitions was experimentally observed in different media [15, 16] when the initial laser pulse duration τ_0 was below 100 fs. In [17–20] it was shown theoretically that, under the conditions of SRSS, the continuous spectral broadening in the Stokes direction occurred when the two-photon Rabi frequency $\Omega^{(2)}$ that corresponded to this stimulated Raman process was less than π/τ_0. But for pulses with more intensity, when the two-photon Rabi frequency was more than π/τ_0 (especially ~$2\pi/\tau_0$), the superbroadening in both the anti-Stokes and Stokes directions and the generation of a spectral supercontinuum occurred [18–20]. In addition, under some conditions, a two-photon self-induced transparency (SIT) at SRSS for few-cycle laser pulses was predicted and an equation for the change in the total energy was derived [17]. However, the transient pulse behavior was not studied. Supercontinuum generation in noble gases by fs pulses was demonstrated for the first time by Corkum *et al.* [21]. Nishioka *et al.* [22] observed white light, ranging from IR (about 1110 nm) to 150 nm, from a self-trapped femtosecond Ti:Al_2O_3 laser pulse (125 fs, 790 nm, 2 TW) in atmospheric-pressure Ar. This was the largest frequency coverage continuum among the rare gases. It is necessary to stress that, at a laser inten-

sity of about 200 TW/cm^2 in argon, the tunneling single ionization time is less than the optical period (the Keldysh parameter equals 0.8), and the ratio of double to single ionization counts Ar^{++}/Ar^+ is less than 0.1 up to a pulse intensity of 1000 TW/cm^2 [23]. Argon is fully ionized at a laser intensity of about 770 TW/cm^2 [24], and the population of Ar neutrals is zero.

Here, we report results of numerical simulations on the resonance generation and propagation of single as well as double attosecond pulses with TW intensity in Ar^+ at SRSS by fs pulses. Two-photon Raman-type self-induced transparency (RSIT) of the pulses is studied. In addition, the transient pulse behavior when input pulse is not in the RSIT regime is investigated to optimize the transition from an input few-cycle pulse to an output attosecond pulse that is in the RSIT regime. The spectral dynamics of the pulses is studied as well. The coherent lengths for plasma mismatching are also calculated for these attosecond pulses.

2. BASIC EQUATIONS

In the Ar^+ ion, the ground state $^2P^o_{3/2}$ and the neighboring excited state $^2P^o_{1/2}$ form a two-level system with the frequency transition $\omega_{21}/2\pi c = 1432$ cm^{-1}, and all other levels are removed on $\Delta \geq 108723$ cm^{-1} [6].

Laser pulses with durations of $\tau = 8$ fs (and shorter) and that have spectral widths of $\delta\nu = 1/(c\tau) = 4167$ cm^{-1} (and more) satisfy the SRSS condition [13, 15]

$$\delta \equiv \tau\omega_{21}/2\pi < 1, \quad (1)$$

where $\hbar\omega_{21} = \epsilon_2 - \epsilon_1$ is the difference in level energies corresponding to the states $^2P^o_{1/2}$ and $^2P^o_{3/2}$, as shown in Fig. 1. The laser pulse field E is described by the wave equation

$$\frac{\partial^2 E}{\partial z^2} - \frac{1}{v^2}\frac{\partial^2 E}{\partial t^2} = \frac{4\pi}{c^2}\frac{\partial^2 P_{nl}}{\partial t^2}, \quad (2)$$

where v is the speed of the laser pulse arising from the linear contribution to polarization and P_{nl} is the nonlinear contribution to polarization caused by the stimulated Raman process. Direct one-photon transitions between the levels $^2P^o_{3/2} - {}^2P^o_{1/2}$ are forbidden by parity in the electronic dipole approximation but permitted for two-photon processes [25] like SRSS. The frequency of Raman transition ω_{21} is considered to be much less than the central frequency ω_c of the initial laser pulse field E. For example, in our case for an initial laser pulse with a wavelength of $\lambda = 800$ nm, the ratio is $\omega_{21}/\omega_c \simeq 0.1$. The condition of strong two-photon Raman-type resonance is completely satisfied by Eq. (1) for all frequency components ω_1 and ω_2 of the laser pulse when, for example, $\delta \leq 0.5$. The dipole polarizability α in rare gases at two-photon resonance is increased by about three orders of

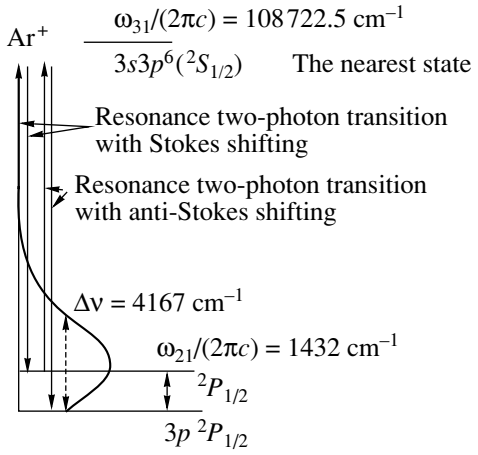

Fig. 1. A diagram showing how the SRSS condition is fulfilled for an 8-fs laser pulse interacting with an Ar^+ ion initially in the ground state $3p\ ^2P^o_{3/2}$. The spectrum (Gaussian) of the pulse, which has a width $\delta\nu = 1/(c\tau) = 4167$ cm^{-1}, contains the frequency transition $\omega_{21}/2\pi c = 1432$ cm^{-1}.

magnitude (experimental results show that the susceptibility χ in rare gases at two-photon resonance can be increased by up to seven orders of magnitude [25]; $\chi \sim |\alpha_{12}|^2$, where α_{12} is a matrix element of the polarizability corresponding to the Raman transition with frequency ω_{21}). Thus, for laser pulses with durations of $\tau = 8$ fs and shorter ($\delta \leq 0.34$), we will have a good approximation (to within one or, at least, a few percents of interaction energy), $V \simeq -\alpha E^2/2$ (where $\alpha = \alpha(\omega_2 - \omega_1 = \omega_{21})$), of the two-photon interaction Hamiltonian [26, 27]. At room temperature and atomic density $N \leq 3 \times 10^{19}$ cm^{-3}, the Doppler dephasing time T^*_D, the polarization relaxation time T_2, and the time of population relaxation T_1 are estimated from experimental data to be $T^*_D \simeq 100$ ps, $T_2 > 0.2$ ps, and $T_2 \ll T_1$, respectively [28, 29]. Hence, first, the condition $\tau < T_2$ of coherent interaction between the pulse and the two-photon transition is very well fulfilled. Second, the inhomogeneous broadening of the resonance transition can be neglected ($\tau/T^*_D < 10^{-3}$). The dynamical Stark shift $\Delta\omega_{st} = (\alpha_{11} - \alpha_{22})E^2/2\hbar$ (α_{11}, α_{22} are values of the diagonal matrix elements of the polarizability) of the resonant levels Kr $6p[5/2]_2$ and the nearest upper $6p[3/2]_1$ and lower $6p[1/2]_1$ (about 200 cm^{-1}) was estimated as $\Delta\omega_{st} \leq 40$ cm^{-1} for an intensity $I = 1$ TW by Kittelmann *et al.* in their work [29] on direct observation of the two-photon coherent response in Kr by femtosecond UV laser pulses. It is necessary to note, first, that the P state in Ar has a dipole polarizability value that is approximately 7.13 times less than that of the P state in Kr [6]. Second, the value of the dipole polar-

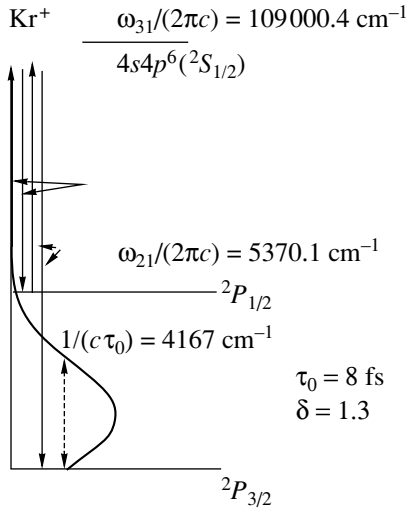

Fig. 2. A diagram showing that the SRSS condition is not fulfilled for an 8-fs laser pulse interacting with a Kr^+ ion initially in the ground state $4p\ ^2P^o_{3/2}$. The spectrum (Gaussian) of the pulse, which has a width $\delta\nu = 1/(c\tau) = 4167\ cm^{-1}$, does not contain the frequency transition $\omega_{21}/2\pi c = 5370\ cm^{-1}$.

izability decreases very rapidly with decreasing principal quantum number. Third, the values of polarizabilities of singly ionized rare gases (in the ground state) heavier than He are approximately two times smaller than for their neutrals [30]. Taking this into account, we can, in our case, neglect the dynamical Stark shift that is connected to the diagonal matrix elements of the polarizability in Ar^+ (for the $3p^5$–$^2P^o_{3/2}$ and $3p(^2P^o_{1/2})$ states), at least for intensities $I \le 10\ TW/cm^2$ ($\Delta\omega_{st} < \omega_{21}$). The nonlinear dipole moment related to SRSS is (see, e.g., [31]) $d = \partial V/\partial E = \alpha E$. Consequently, the nonlinear contribution to the polarization of the medium included in Eq. (2) is $P_{nl} = N\langle\alpha\rangle E$, where N is the number density of Ar^+ ions and an average polarizability value $\langle\alpha\rangle = \mathrm{Sp}(\alpha\rho)$ can be derived by using the density matrix ρ. The function $\langle\alpha'\rangle = \langle\alpha\rangle/|\alpha_{12}|$ and the difference between the level populations $w = \rho_{22} - \rho_{11}$ are described by the equations [13, 27]

$$\frac{\partial^2\langle\alpha'\rangle}{\partial t^2} + \frac{2}{T_2}\frac{\partial\langle\alpha'\rangle}{\partial t} + \omega_{21}^2\langle\alpha'\rangle = -\omega_{21}\Omega^{(2)}w, \qquad (3)$$

$$\frac{\partial w}{\partial t} + \frac{(w - w_0)}{T_1} = \frac{\Omega^{(2)}}{\omega_{21}}\left(\frac{\partial\langle\alpha'\rangle}{\partial t} + \frac{\langle\alpha'\rangle}{T_2}\right), \qquad (4)$$

where $\Omega^{(2)} = |\alpha_{12}|E^2/\hbar$ is the two-photon Rabi frequency. For further consideration, it is convenient to introduce an angle of rotation of the material variable [17–19]:

$$\varphi(t, z) = \frac{|\alpha_{12}|}{\hbar}\int_{-\infty}^{t} E^2(t', z)dt'. \qquad (5)$$

The quantity $S(z) = \varphi(+\infty, z)$ is the integral over the entire pulse of the two-photon Rabi frequency of the process. It is easy to see that $S(z)$ is in direct proportion to the total energy of the laser pulse that has passed through a distance z. The value of the nonresonant dipole polarizability for Ar in the S state is equal to about $1.64 \times 10^{-24}\ cm^3$, and, in the P state, it is equal to about $7.1 \times 10^{-24}\ cm^3$ [6]. As was noted above, values of polarizabilities of singly ionized rare gases (in the ground state) heavier than He are approximately two times less than of their neutrals. For example, the ion Ne^+ and the atom Ne have values of static polarizabilities of $0.21 \times 10^{-24}\ cm^3$ and $0.395 \times 10^{-24}\ cm^3$, respectively [30]. Thus, with a knowledge of the dynamic polarizability for Ar at 800 nm, and taking into account its above-mentioned increase at the two-photon resonance, the value of the matrix element of the polarizability $|\alpha_{12}|$ (corresponding to the transition $^2P^o_{3/2}$–$^2P^o_{1/2}$ in Ar^+) can be estimated to be in the range 10^{-21}–$10^{-20}\ cm^3$. With regard to Kr, the value of its static polarizability ($2.48 \times 10^{-24}\ cm^3$ [30]) is larger but still comparable to the Ar value. However, the fine structure $^2P^o_{3/2}$–$^2P^o_{1/2}$ levels of Kr^+ ions (where $^2P^o_{3/2}$ is the ground state) form a two-level system with the frequency transition $\omega_{21}/2\pi c = 5370\ cm^{-1}$, and all other levels are removed on $\Delta > 112828\ cm^{-1}$ [6], as is shown in Fig. 2. Thus, to satisfy SRSS condition (1) for Kr^+, it is necessary to have a laser pulse with a duration of $\tau \le 6$ fs.

The numerical solutions of wave equation (2) with Eqs. (3) and (4) in the case of a coherent interaction for a laser pulse when $S(0) = 2\pi n$ ($n = 1, 2, 3$) and for various $\beta z \le 2$, where $\beta = 2\pi N|w_0||\alpha_{12}|\omega_{21}/c$ ($N \le 3 \times 10^{19}\ cm^{-3}$), have shown that no more than one percent (within 1% accuracy of the simulation) of the pulse energy was scattered in the reverse direction [13]. (The quantity β^{-1}, which has the dimension of length, can be interpreted as the length of coherent SRSS.) It is therefore a very good approximation to introduce a reduced wave equation [13, 17–20]:

$$\frac{\partial E}{\partial z} + \frac{1}{v}\frac{\partial E}{\partial t} = -\frac{4\pi v}{c^2}\frac{\partial P_{nl}}{\partial t}, \qquad (6)$$

which describes only the forward scattering of the field. First, the reduced wave equation was derived by Bullough, Eilbek *et al.* for one-photon SIT [32]. If a dimensionless constant $\alpha_s^{(1)} = 4\pi N d_{21}^2(\hbar\omega_{21})^{-1}$ (where N is an atomic density and d_{21} is a matrix element of the

one-photon resonant transition with frequency ω_{21}) is less than 1, then the reduced Maxwell–Bloch equations (the Maxwell–Bloch equations in which only backscattering is neglected) provides a good description of SIT for pulses with durations of up to a half cycle of the electromagnetic field. For example, for a typical value of d_{21}, namely, 10^{-18} cgs units, this is a good approximation (to within 1%) at $N = 10^{18}$ cm^{-3}. It should be stressed that the numerical simulations using Eq. (6) use significantly less computer time than Eq. (2). In addition, when $\delta \ll 1$, the laser pulse energy is altered by the law [17, 19]

$$\frac{dS}{dz} = \beta(1 - \cos S), \qquad (7)$$

which is practically the same as for the case of a coherent direct two-photon interaction of a field with a resonant medium, which was first derived by Belenov and Poluektov [33]. This is the area theorem for two-photon self-induced transparency. Such laser pulses with $S(0) = 2\pi n$ can propagate through a medium ($\beta z \approx 1$) in the RSIT regime [17–20]. Furthermore, the propagation of the pulses with $S = 2\pi$ [12, 13, 18, 20], $S = 4\pi$ [13], and $S = 6\pi$ [34] leads to effective self-broadening of the spectrum and to the generation of a spectral supercontinuum due to SRSS.

3. RESULTS AND DISCUSSION

The dynamics of a laser pulse having an initial Gaussian shape given by

$$E(t, 0) = E_0 \sin(\omega_c t) \exp\left\{-2\ln 2\left(\frac{t}{\tau_0}\right)^2\right\}, \qquad (8)$$

where E_0 is the initial peak field amplitude of the pulse propagating through the initially unexcited medium (i.e., ($w(0, z) = w_0 = -1$, $\langle\alpha(0, z)\rangle = 0$)), is investigated numerically by Eqs. (3), (4), and (6). An improved and repeatedly tested version of the code presented in [12, 13, 17–20] is used on a grid of 4096 points per one period of an electromagnetic field. This provided for a good solution (better than 1% accuracy). As far as we know, there are no direct experimental data on the dynamic polarizability or the two-photon Rabi frequency for the two-photon Raman-type resonance at the transition $^2P^o_{3/2} - {}^2P^o_{1/2}$ in Ar^+ that could be used to estimate directly the value of the matrix element of the polarizability $|\alpha_{12}|$. Taking into account the above estimations of $|\alpha_{12}|$, we performed numerical simulations for a broad range of the parameter $\gamma = 2\pi N|\alpha_{12}| \simeq 0.04$–0.2, and we found similar results as some function of βz.

In Fig. 3 we show the dynamics of the field E of an initial three-cycle (within FWHM) laser pulse ($\beta z = 0$) (Fig. 3a) when $\gamma = 0.2$, $S(0) = 6\pi$, and its field at $\beta z = 0.34$ (Fig. 3b), $\beta z = 0.51$ (Fig. 3c), and $\beta z = 0.68$ (Fig. 3d). The field E that has the double half-cycle pulse of Fig. 3d is displayed in more detail in Fig. 4. The dynamics of its spectrum is illustrated in Fig. 5, in Fig. 5a at $\beta z = 0$, and in Fig. 5b at $\beta z = 0.68$. This 6π pulse propagated through a distance corresponding to

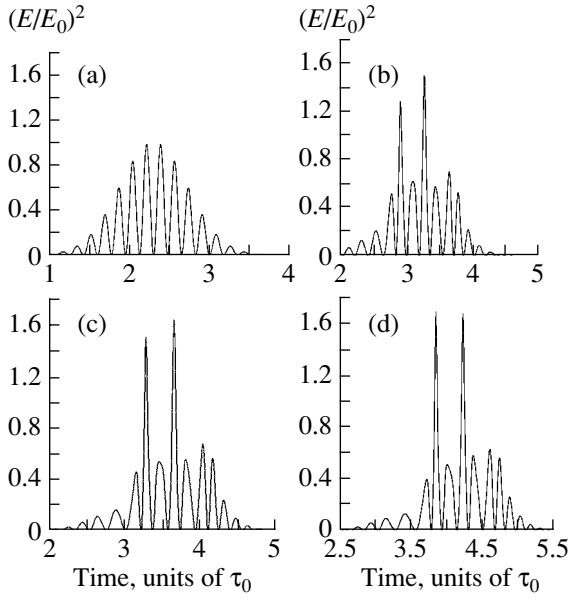

Fig. 3. Dynamics of the squared field E (normalized to the initial amplitude E_0) vs. time (normalized to the initial pulse width τ_0) for an initial three-cycle (FWHM) laser pulse in Ar^+: (a) ($\beta z = 0$), $S(0) = 6\pi$, $\gamma = 0.2$, at (b) $\beta z = 0.34$, (c) $\beta z = 0.51$, and (d) $\beta z = 0.68$.

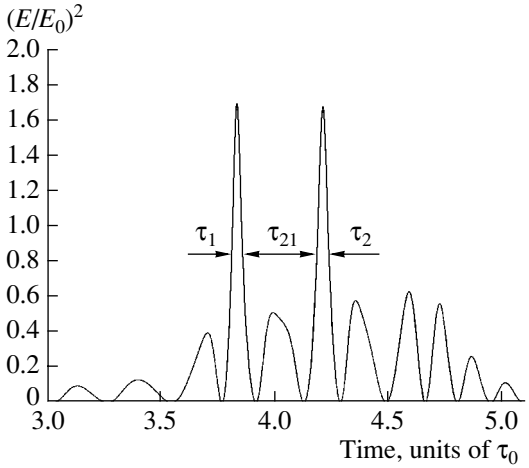

Fig. 4. The field E with the double half-cycle pulse of Fig. 3d is displayed in more detail. The first and second 2π pulses are transformed into half-cycle (FWHM) pulses with greater amplitudes and photon energies. The field lengths of the 2π pulses are τ_1 and τ_2 (FWHM), and τ_{21} is the time distance between these pulses. For $\beta z = 0.68$, $\tau_1 = 0.14T_0$, $\tau_2 = 0.15T_0$, and $\tau_{21} = 1.12T_0$. (T_0 is a period of the oscillation of the initial field E.)

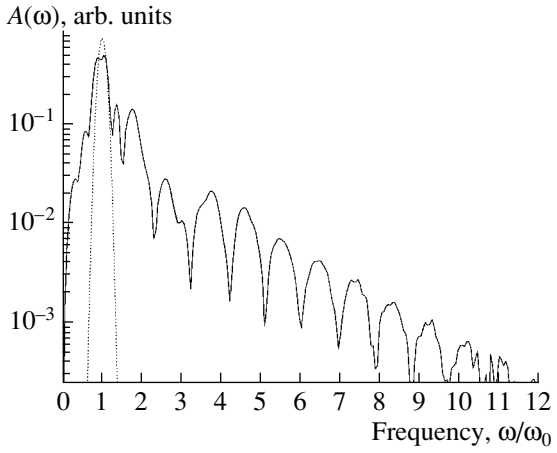

Fig. 5. Amplitude spectra $A(\omega)$ of the field E, corresponding to Figs. 3a and 3d. The dotted line is the spectrum of the initial field ($\beta z = 0$, Fig. 3a). The solid line is the spectrum of the field after a distance corresponding to $\beta z = 0.68$ (Fig. 3d).

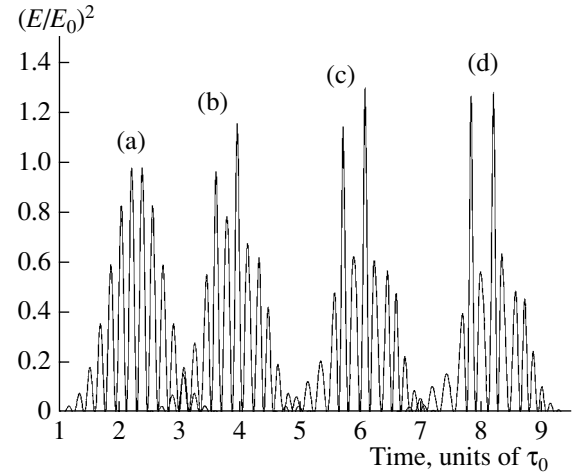

Fig. 6. Dynamics of squared field E (normalized to initial amplitude E_0) vs. time (normalized to initial pulse width τ_0) for an initial three-cycle (FWHM) laser pulse in Ar^+: (a) ($\beta z = 0$), $S(0) = 6\pi$, $\gamma = 0.04$, at (b) $\beta z = 0.18$, (c) $\beta z = 0.36$, and (d) $\beta z = 0.54$.

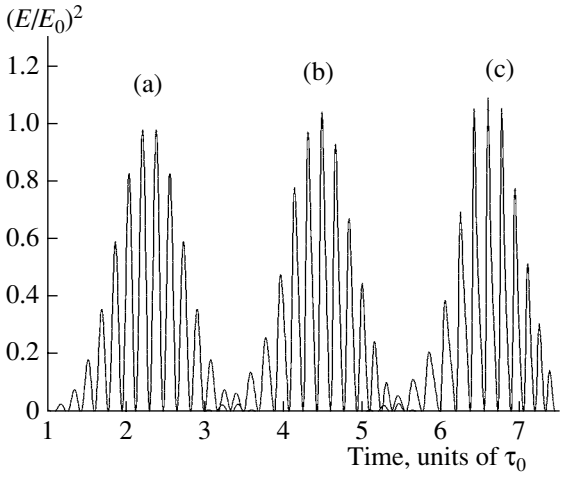

Fig. 7. Dynamics of the squared field E (normalized to initial amplitude E_0) vs. time (normalized to initial pulse width τ_0) for an initial three-cycle (FWHM) laser pulse in Kr^+: (a) ($\beta z = 0$), $S(0) = 6\pi$, $\gamma = 0.04$, at (b) $\beta z = 0.64$, and (c) $\beta z = 1.28$.

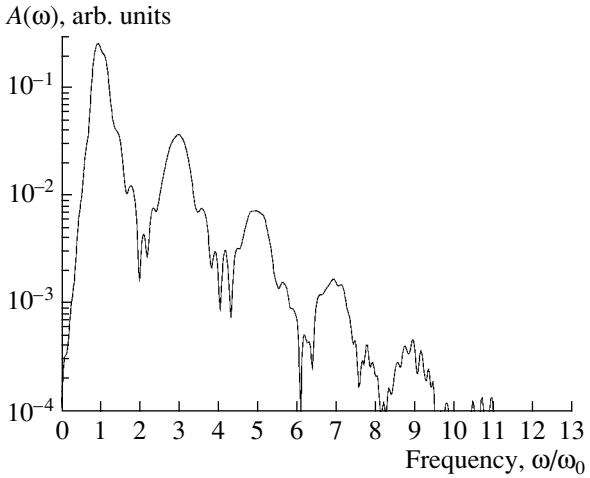

Fig. 8. Amplitude spectrum $A(\omega)$ of the field E, corresponding to Fig. 7c, after a propagation distance corresponding to $\beta z = 1.28$ (Fig. 7c).

$\beta z = 0.68$ in the RSIT regime while conserving about 98% of its energy (to within 1%). The initial 6π pulse split into three 2π pulses (Fig. 3d) that, with the conserved total energy of the initial pulse, are unique signatures of self-induced transparency. Some properties of the two-photon Raman-type self-induced transparency for few-cycle laser pulses were studied recently [34]. In particular, it was shown that, after such splitting of an initial six-cycle (FWHM) 6π pulse into three 2π pulses, the first and second 2π pulses were transformed into one-cycle (FWHM) pulses with greater amplitudes and photon energies (see Figs. 5 and 6 of [34]). Figure 3 shows, in the case of an initial three-cycle (FWHM) 6π pulse, the first and second 2π pulses are transformed into half-cycle (FWHM) pulses with greater amplitudes and photon energies [35]. The field lengths of the 2π pulses are decreased and, at $\beta z = 0.68$, are $\tau_1 = 0.14T_0$ and $\tau_2 = 0.15T_0$ (FWHM), where T_0 is the period of one oscillation of the initial field E ($T_0 = 2\pi/\omega_c$). The time distance between these pulses is $\tau_{21} = 1.12T_0$ (FWHM). For example, if an initial three-cycle laser pulse has a wavelength $\lambda = 800$ nm, then $T_0 = 2.667$ fs, $\tau_0 = 8$ fs, $\tau_1 = 0.37$ fs, $\tau_2 = 0.40$ fs, and $\tau_{21} = 2.99$ fs. The dynam-

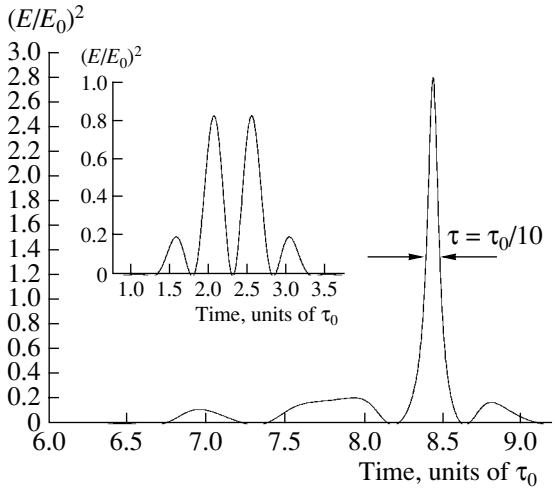

Fig. 9. Dynamics of the squared field E (normalized to initial amplitude E_0) vs. time (normalized to initial pulse width τ_0) for an initial one-cycle (FWHM) laser pulse shown in inset: (a) $S(0) = 2\pi$, $\gamma = 0.2$, and (b) $\beta z = 0.78$.

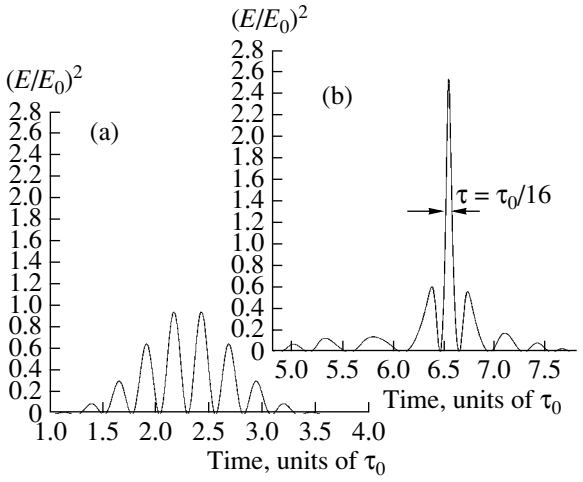

Fig. 10. Dynamics of the squared field E (normalized to initial amplitude E_0) vs. time (normalized to initial pulse width τ_0) for an initial two-cycle (FWHM) laser pulse: (a) ($\beta z = 0$), $S(0) = 2.7\pi$, $\gamma = 0.2$, and (b) $\beta z = 1.5$.

ics of the amplitude spectrum $A(\omega)$ of the field E corresponding to Fig. 3 is shown in Fig. 5.

In Fig. 6 we show the dynamics of the field E of the same initial three-cycle laser pulse ($\beta z = 0$) (Fig. 6a), but for when $\gamma = 0.04$. We found results similar to Fig. 3, but for other βz and corresponding propagation distances.

To illustrate the essence of the SRSS condition for this effect, we show the dynamics of the field E in Fig. 7, as well as the spectra in Fig. 8, for the same initial three-cycle laser pulse in Kr^+ ions that are initially in the ground state $4p\ ^2P^o_{3/2}$. It is easy to see (Fig. 7) that, in this case, we have no splitting into three 2π pulses and double half-cycle (FWHM) pulses with greater amplitudes. The spectrum is of odd harmonics, as can be seen in Fig. 8. Thus, due to the discrete spectra, only trains of attosecond pulses with the abovementioned small efficiency can be generated in this case.

Next, we turn to the Ar^+ ions that are initially in the ground state. One of the best results was found for a one-cycle pulse (FWHM) (Fig. 9, inset) at $S(0) = 2\pi$, $\gamma = 0.2$. This pulse propagated under RSIT conditions, transforming into a single half-cycle pulse with a duration $\tau = \tau_0/10 = 0.1T_0$ at $\beta z = 0.78$, as is shown in Fig. 9. For example, if an initial one-cycle laser pulse has a wavelength $\lambda = 800$ nm, then $T_0 = 2.7$ fs, $\tau_0 = 2.7$, and $\tau = 0.27$ fs.

Because, at present, the experimental technique to generate two-cycle (FWHM) laser pulses with TW intensity is still being developed [4], we investigated the feasibility of generating a single attosecond pulse from such two-cycle pulses. Our numerical simulation shows that some of the best results can be obtained if initial two-cycle laser pulses with $S(0) = 2.7\pi$ are used [35]. Thus, the pulses do not propagate initially in the RSIT regime. However, after some distance ($\beta z \simeq 1.5$), they will have $S(z) \simeq 2\pi$ and will propagate under RSIT conditions. The results with the two-cycle (FWHM) input pulse are presented in Fig. 10, and the corresponding spectra are shown in Fig. 11. Passing through some distance corresponding to $\beta z \simeq 1.5$, the two-cycle pulse (Fig. 10a) is transformed into a half-cycle pulse (FWHM) (Fig. 10b) with the duration $\tau = \tau_0/16 = 0.125T_0$. The spectrum of its field is a smooth supercontinuum (Fig. 11), as compared with one of the double half-cycle pulses (Fig. 3). For example, if an initial two-

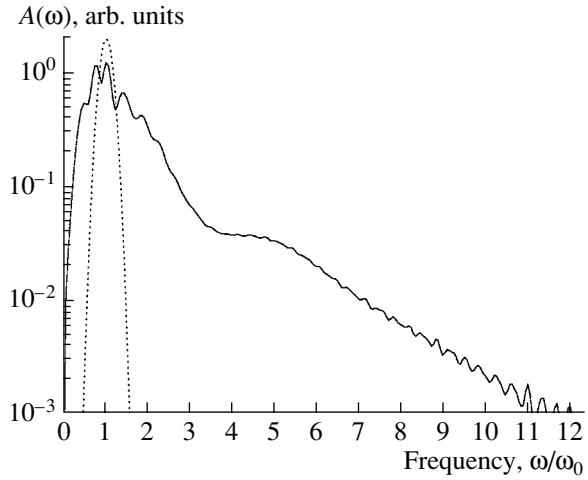

Fig. 11. Amplitude spectra $A(\omega)$ of the field E corresponding to Fig. 10. The dotted line is the spectrum of the initial field ($\beta z = 0$). The solid line is the spectrum of the field after a distance corresponding to $\beta z = 1.5$.

cycle laser pulse has a wavelength $\lambda = 800$ nm, then $T_0 = 2.67$ fs, $\tau_0 = 5.34$, and $\tau = 0.33$ fs.

Using the aforementioned parameters of the numerical simulations, we can give as an example some corresponding concrete parameters for future experimental investigations. It is easy to see that $\gamma \simeq 0.2$ and $\beta \simeq 1.8 \times 10^3$ cm^{-1} if, for example, $N = 3 \times 10^{18}$ cm^{-3} and $|\alpha_{12}| = 10^{-20}$ cm^3 or $N = 3 \times 10^{19}$ cm^{-3} and $|\alpha_{12}| = 10^{-21}$ cm^3. For a 5-fs two-cycle laser pulse and $|\alpha_{12}| = 3 \times 10^{-21}$ cm^3, we will have $S(0) \simeq 3\pi$ at $I \simeq 10^{12}$ W/cm^2, and the propagation distances $z \simeq 8$ μm as in Fig. 10b. In addition, using the same method as in [34], we have calculated the effect of the plasma on coherent lengths. For example, at densities $N = 10^{19}$ cm^{-3} of Ar^+ ions and free electrons, as well as for the initial 5-fs two-cycle laser pulse with $\lambda = 800$ nm, this coherent length $L_c \simeq 12$ μm for the 330-as pulse (Fig. 9b) that has a spectral supercontinuum corresponding to Fig. 11.

4. SUMMARY

In summary, we have shown that three-, two-, and one-cycle laser pulses with lengths of less than 10 fs can be transformed efficiently in Ar^+ into double and single half-cycle attosecond pulses, respectively, by resonant two-photon Raman-type transitions. In addition, these pulses, with TW-range peak intensities, can propagate through some distance in the regime of two-photon Raman-type self-induced transparency. The coherent lengths for plasma mismatching have been calculated and are longer than the generation lengths of the attosecond pulses.

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