Shaping of few-cycle laser pulses via a subwavelength structure*

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We theoretically investigate the propagation of few-cycle laser pulses in resonant two-level dense media with a subwavelength structure, which is described by the full Maxwell–Bloch equations without the frame of slowly varying envelope and rotating wave approximations. The input pulses can be shaped into shorter ones with a single or less than one optical cycle. The effect of the parameters of the subwavelength structure and laser pulses is studied. Our study shows that the media with a subwavelength structure can significantly shape the few-cycle pulses into a subcycle pulse, even for the case of chirp pulses as input fields. This suggests that such subwavelength structures have potential application in the shaping of few-cycle laser pulses.

Keywords: ultrashort pulses, coherent propagation, self-induced transparency

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1. Introduction

With the development of ultrafast science and technology, the generation of light pulses with durations down to a few optical cycles has been permitted.[1] As we know, electrons play a key role in the physical, chemical, and biological processes in the microcosm. Ultrashort and ultraintense few-cycle laser pulses can be used to study the electronic dynamics on the atomic and molecular scales.[1,2] Although many techniques have been explored to obtain few-cycle pulses, the generation of optical pulses with one or fewer currently is one of the important targets in the applications of the ultrafast laser. The light–matter interaction under extreme conditions is significant as the pulse approaches the duration of single or less than one optical cycle.[3–8] For example, the spectral width of such ultrashort pulses becomes comparable to the carrier frequency, which is referred to spectral supercontinuum.[9] For the light–matter interactions in such a case, the traditional frame of slowly varying envelope approximation (SVEA) and rotating wave approximation (RWA) is invalid.[10] Many interesting propagation phenomena have been found, for example, the formation of few- and single-cycle solitons,[3–5,11] self-induced transparency (SIT),[8,12] supercontinuum radiation,[13,14] and carrier-wave Rabi flopping.[15] With an in-depth study of the propagation of the few-cycle pulses, it has been found that it is significant to design a device to shape few-cycle pulses and get shorter pulses.[1,16–21]

The theory and experiment of periodic optical structures, such as the Bragg grating, photonic band-gap crystals, and waveguide arrays have achieved a rapid progress in the last decade.[22] Such systems can be used to control the light propagation and light–matter interaction and have many applications, in particular, such as nonlinear filtering, switching, and distributed-feedback amplification.[23] In the case of few-cycle pulses penetrating through continuous media, big blueshifts occurred,[8] which are caused by the intrapulse four-wave mixing inducted by the rotating term. Such a property limits the applications of the continuous media in the shaping of ultrashort pulses. Many interesting phenomena have been revealed in periodic structures, such as the gap-soliton generation,[24–27] whose mechanism is essentially different from that in continuous media. Periodic structures usually consist of a periodic array of thin layers of a resonant two-level medium separated by half-wavelength nonabsorbing dielectric layers (Bragg condition). Recently, it was reported that a single-cycle gap soliton could be supported by a subwavelength structure in which the thickness of each two-level medium layer was much less than the wavelength of the electromagnetic wave propagating.[28] At the same time, one scheme to generate an ultrashort pulse with periodic subwavelength structure has been reported.[29] In this paper, we will investigate the potential application of subwavelength structures in the ultrashort-pulse shaping. Many mathematical models have been established to research the ultrashort pulse propagation from different points of view, such as the full Maxwell–Bloch model,[11,28,30] the nonlinear Schrödinger equation with higher order corrections,[6] the modified Korteweg-de Vries and sine-Gordon equations,[3,4,31,32] and so on. It is worth noting that the last models may be derived from the full Maxwell–

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Bloch model. In order to describe the interaction of light pulses and the two-level system embedded in a Bragg grating, we take a semiclassical Maxwell–Bloch model without the frame of slowly varying envelope and rotating wave approximations as in Ref. [28].

The content of the paper is organized as follows: we first show our model to describe the few-cycle pulse propagation in a periodic structure which consists of two-level medium in Section 2; In Section 3, several numerical results in different kinds of conditions are presented, meanwhile, brief discussions are given respectively. Finally, we present a brief summation.

2. Model

In order to study the shaping effect in the media structure, we will consider the following cases. Let an extremely short pulse propagate along ξ (or z) axis in vacuum, then it arrives at an input interface of a one-dimensional periodic structure at $\xi = 0$ (a set of dimensionless time and space coordinates were taken, for example, $T = t\omega_0$ and $\xi = z/\lambda_0$ with the carrier frequency ω_0 and wavelength λ_0). At first, the few-cycle pulse is propagating in the vacuum. Then the pulse partially penetrates into the subwavelength structure. Finally, the penetrating part passes through and exits into the free space. The periodic structure consist of periodic thin layers in which resonant two-level media is contained and separated by a transparent material is considered (see in Fig. 1(a)), which differs from the ordinary structures of a Bragg grating. For normal atomic and molecular systems, the intensities of ultrashort pulses with Rabi frequency $\Omega \simeq \omega$ (the transition frequency) is so high that the ionization phenomena should be considered because the corresponding dipole momentums are too small. In order to carry out the experiment, some artificial atomic systems, such as quantum dots[33] and quantum wells,[34] can be applied to investigate the dynamics in strong-coupling regime. One alternative experiment system is symmetric coupled double quantum wells consisting of two 5.09-nm-wide GaAs wells, separated by a 1.13-nm-wide pure AlAs central barrier, and sandwiched between $Al_{0.3}Ga_{0.7}As$ layers.[35] In such a system, there are only four bound states, as shown in Fig. 1(b). When the bias voltage is zero and the temperature is 45 K, the energy of the interband transition is about 1620 meV (766.9 nm), the energy gap Δ of the two excited levels is about 1 meV.[35] The dipole matrix element of the corresponding electron–hole transitions is eD (D is the interwell distance and is about 3.5 nm).[36] If $\Omega = \omega$, we can get the corresponding electric field amplitude being $\sim 4.63 \times 10^8$ V/m. According to the formulae $E = 2.74 \times 10^3\sqrt{I}$, where E and I are in units of V/m and W/cm^2, respectively, the peak intensity is $\sim 2.85 \times 10^{10}$ W/cm^2. In such intensity regime, the ionization can be neglected. At the same time, the energy gap Δ between $|1\rangle \to |2\rangle$ is much smaller than that between the transition $|0\rangle \to |1\rangle$, the two excited states can be served as one state because the duration of the light-matter interaction is much less than the time scale $1/\Delta$. Thus, we can conclude that a practical two-level model is obtained to study a quantum system interacting with a ultrashort pulse in some conditions. The evolutions of the few-cycle electromagnetic field amplitudes $E_x(T,\xi)$ and $H_y(T,\xi)$ are described by Maxwell's equations which take the form

$$\begin{aligned} \frac{\partial}{\partial T}H &= -\frac{1}{2\pi}\frac{\partial}{\partial \xi}E, \\ \frac{\partial}{\partial T}E &= -\frac{1}{2\pi}\frac{\partial}{\partial \xi}H - \frac{1}{\varepsilon_0 E_0}\frac{\partial}{\partial T}P. \end{aligned} \tag{1}$$

Here, $E = E_x/E_0$ and $H = \sqrt{\mu_0/\varepsilon_0}H_y/E_0$ are the normalized dimensionless components of the few-cycle pulse with the peak amplitude of the electric component of the incident pulse E_0. The atomic polarization $P = N(\xi)gu$ is determined by the off-diagonal density matrix element $\rho_{01} = u(T,\xi)/2 + \mathrm{i}v(T,\xi)/2$, $N(\xi)$ and g denote the density distributions of active medium and the dipole moment, respectively. Polarization variables u, v, and inversion population $w(T,\xi) = \rho_{11} - \rho_{00}$ obey the following set of Bloch equations

$$\begin{aligned} \frac{\partial}{\partial T}u &= -\frac{1}{\tilde{T}_2\omega_0}u - \frac{\omega}{\omega_0}v, \\ \frac{\partial}{\partial T}v &= -\frac{1}{\tilde{T}_2\omega_0}v + \frac{\omega}{\omega_0}u + \frac{2\Omega_0}{\omega_0}E\,w, \\ \frac{\partial}{\partial T}w &= -\frac{1}{\tilde{T}_1\omega_0}(w+1) - \frac{2\Omega_0}{\omega_0}E\,v. \end{aligned} \tag{2}$$

Here, $\Omega_0 = gE_0/\hbar$ is the peak Rabi frequency of the incident pulse; $\tilde{T}_1$ and $\tilde{T}_2$ are, respectively, the lifetimes of the excited state and the dephasing time. For analytical convenience, we introduce new parameter $\omega_c(\xi) = N(\xi)g^2/\varepsilon_0\hbar$ to denote the density distribution, which is given by the expression

$$\omega_c(\xi) = \begin{cases} d\ \mathrm{fs}^{-1} & \text{if } \xi \in [2N_s\delta - 2\delta,\ 2N_s\delta - \delta), \\ 0 & \text{if } \xi \in [2N_s\delta - \delta,\ 2N_s\delta), \end{cases} \tag{3}$$

where d and δ express the density and the thickness of one layer; N_s denotes the number of the layers consisted with resonant two-level medium and is an integer with rang $[1,\ L/(2\delta)]$. The length of the periodic structure is L. The density distribution of atomic medium is shown in Fig. 1(a). Using parameter $\omega_c(\xi)$ and substituting the equation for u into Maxwell's equation (1), we then obtain

$$\begin{aligned} \frac{\partial}{\partial T}H &= -\frac{1}{2\pi}\frac{\partial}{\partial \xi}E, \\ \frac{\partial}{\partial T}E &= -\frac{1}{2\pi}\frac{\partial}{\partial \xi}H - \frac{\omega_c(\xi)}{\Omega_0}\frac{\partial}{\partial T}u. \end{aligned} \tag{4}$$

Equations (2) and (4) constitute a theoretical model of our system, which is named as the full Maxwell–Bloch equations.

The incident pulse we consider is a few-cycle pulse with a hyperbolic secant envelope, i.e., $\Omega(T=0,\xi)=\Omega_0\cos[2\pi(\xi-\xi_0)+\varphi_0-2\pi\alpha(\xi-\xi_0)^2]\mathrm{sech}[1.76\times 2\pi(\xi-\xi_0)/\tau_\mathrm{p}]$, where α is the chirp coefficient and τ_p is the full width half maximum (FWHM) of the pulse intensity envelope, ξ_0 and φ_0 labels the position and the carrier envelope phase of the initial pulse's peak, respectively. The two-level medium is initialized with $u=v=0$ and $w=-1$ at $t=0$. The material and laser parameters are chosen as follows: $\tau_\mathrm{p}=4.69$ fs, $\varphi_0=0$, $\omega_0=2.45$ fs^{-1} ($\lambda_0=766.9$ nm), $g=5.61\times10^{-29}$A$\cdot$s$\cdot$m, $T_1=0.939$ ps, $T_2=0.470$ ps, and $N=5.61\times10^{17}$ cm^{-3}. In addition, we set $\omega=\omega_0$, which corresponds to the resonant coupling The above parameters give the atomic density $d=0.2$. When $\Omega_0=1.49$ fs^{-1} and $\alpha=0$, it is easy to get that the input envelope area is $A=\int\Omega(t)\,\mathrm{d}t=\Omega_0\tau_\mathrm{p}\pi/1.76=4\pi$ and the corresponding cycle number is $N_\mathrm{c}=\Omega_0\tau_\mathrm{p}/2\pi=1.83$. In all simulations, we take that $\xi_0=18$ in order to ensure that the pulse locates in the free space and the partial pulse in the structure can be neglected at $t=0$.

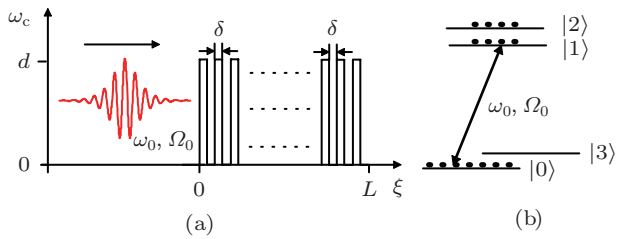


Fig. 1. (color online) (a) Spatial distribution of active medium. The red curve illustrates the incident wave with a few optical cycles, which will propagate through the resonant two-level dense media with a periodic structure. Symbol $\rightarrow$ denotes the propagation direction ξ (or z) axis, ω_0 and Ω_0 denotes the carrier frequency and the Rabi frequency of the initial pulse's peak, respectively. (b) Scheme of the energy levels of symmetric coupled double quantum wells.[35] When a linear polarized laser field is applied, the transitions $|0\rangle\rightarrow|1\rangle$ and $|0\rangle\rightarrow|2\rangle$ are coupled.

3. Numerical results and discussion

The full Maxwell–Bloch equations (2) and (4) for different periodic structures and conditions are solved by using a finite difference time-domain method[30] which has been proven to be an accurate tool. First, we investigate the few-cycle pulse propagation in a resonant structure in which the structure period 2δ exactly satisfies the Bragg condition: $2\delta=m/2$. Here, the length is expressed in the dimensionless coordinate and m is an integer. In our case, the period of the structure is set to be equal to $1/2$ ($\delta=1/4$). In our simulation, the medium and structure parameters are: $L=108\lambda_0$, $d=0.4$, $\tau_\mathrm{p}=4.69$ fs, $\delta=0.25$, $\Omega_0=1.49$ fs^{-1}. In such conditions, the peak amplitude of the electric field is $E_0=2.815\times10^8$ V/m. Similar electric-field profiles of the penetrating pulse inside and outside the periodic structure for different time moments to those in Fig. 2 of Ref. [28] can be gotten. Here, the cycle number of the transmitted pulse is reduced in comparison to the initial pulse and approaches to a ultrashort pulse with a single optical cycle. It can be found that the area of the penetrating pulse is 2π, which agrees well with that in the case that a pulse is propagating in a resonant absorptive medium. This phenomenon is known as the area theorem.[37,38] We can also obtain the evolution of the pulse spectrum during propagation through the structure above. As in Fig. 4 in Ref. [28], we observe that the reflected and the penetrating spectra are situated around the input central frequency ω_0, which is distinct from the results of few-cycle pulse propagation in the continuous media.[8] In the case of continuous media, there is a large frequency shift of the transmitted and reflected pulses caused by the intrapulse four-wave mixing. In our case, however, this shift can be efficiently suppressed. As we know, periodic structures modify the dispersion property of the media and offer optical band gaps. Such modification destroys the phase matching condition of the four-wave mixing process, which guarantees that the large frequency shift does not occur in our simulations. From here, we conclude that the periodic structure provides an alternative method to reshape few-cycle pulses and allows one to obtain single-cycle or subcycle optical pulses.

Subsequently, in order to investigate the effect of the density d on the propagation, we increase the density of the two-level system. Finally, the single-cycle pulse can not be obtained. Here, all parameters are the same as these in the former simulation, however, for different d. This is totally different from the case in the continuous media, in which increasing the densities will lead to even larger frequency shifts of the output pulses (see, for example, Ref. [8]). To show this, the evolutions of the electric field profiles of the propagating pulse for different time moments are depicted in Fig. 2, from which one see that when the initial pulse enters into the dense medium, the energy of the penetrating pulse becomes dislocated and looks random (see Fig. 2(d)). The reason is that the reflection on the interfaces of the excited layers is enhanced rapidly with the increase of the density d. The energy of the penetrating field spreads and becomes weak (weaker than the reflected one). Most of the propagating pulse energy locates in the periodic structure or is reflected. The higher density will also result in an enhanced reflection (see Fig. 2(a)). The reason is that the reflection on the interfaces of the media layers is enhanced rapidly with the increase of the density of the two-level media. The output field spreads and becomes weak. Similar behavior occurs if one moderately increases the initial pulse intensity. However, a shorter structure ($\xi<20$) can also be used to achieve a output pulse with one or less than one cycle (see Fig. 2(a)). If we increase the density continuously, the penetrating pulse will decay more quickly and disappear

after a shorter distance. Hereto, we have discussed how the media density affects the pulse propagation and conclude that the density in a suitable regime is important to shape few-cycle pulses.

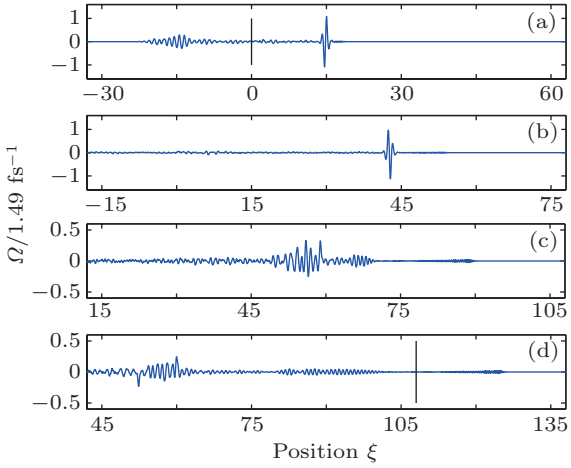


Fig. 2. (color online) Electric-field profiles for different time moments. (a) $t = 0.094$ ps; (b) $t = 0.188$ ps; (c) $t = 0.282$ ps; (d) $t = 0.376$ ps. The other medium and structural parameters are: $\xi \in [0,\ 108]$, $d = 0.4$, $\tau_p = 4.69$ fs, $\delta = 0.25$, $\Omega_0 = 1.49$ fs^{-1} (the corresponding peak amplitude of electric field $E_0 = 2.815 \times 10^8$ V/m). (a) and (d) correspond to the electric-field profiles near the input face and outside the exit face [vertical lines in (a) and (d)], respectively.

Then, we investigate the impact of the chirp coefficient α on the pulse propagating in a resonant structure. Electric field profiles of the penetrating pulse inside and outside the periodic structure for different time moments are shown in Fig. 3. It is easy to see that in the case of $\delta = 0.25$, the impact of the chirp parameter is relatively significant on the output fields, such as the group velocity. If the periodic structure becomes more compact, the impact of the chirp parameter is suppressed (see the left plots in Fig. 4). Here, we see that even in the case where the initial pulse is not ideal, a similar pulse shape can be gotten after propagation when the structure period is compact. This feature is beneficial to the shaping of the arbitrary few-cycle pulses. It should be noticed that the structure we considered here breaks the Bragg condition ($\delta = 0.1$).

In the following, the impact of the increasing pulse area on the few-cycle pulse propagating is studied. A pair of pulses with few optical cycles are observed in our model, which is shown in Fig. 3. According to the SIT theorem, the pulse with the initial envelope area $n \times 2\pi$ will be separated into n subpulses with the envelope area 2π.[37,38] In Fig. 3, the initial envelope area we considered is much bigger than 4π. Only two split pulses exist in our model because a part of the energy of the input laser pulse is reflected and lost. In addition, we observe that the stable pulses only have a few differences in amplitude, velocity, and duration, but have the same pulse area (2π). The corresponding analytical solution is difficult to be derived until now. From the viewpoint of the shaping efficiency, a tighter structure is more effective because more energy can enter into the media [see Figs. 3(b) and 3(d)]. Here, it should be noted that the smaller pulses in Fig. 3(b) disappear when the length of the optical structure is increased. The reason is that all pulses continue to broaden in the media if the envelope area is less than 2π, as mentioned above.

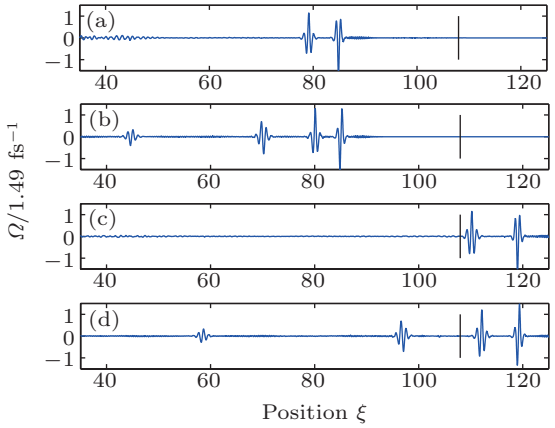


Fig. 3. (color online) Electric-field profiles for different time moments and structure periods near the output interface. Here a longer pulse without chirp is considered ($\tau_p = 11.7$ fs and the corresponding pulse area is 10π). (a) $t = 0.282$ ps, $\delta = 0.25$; (b) $t = 0.282$ ps, $\delta = 0.1$; (c) $t = 0.376$ ps, $\delta = 0.25$; (d) $t = 0.376$ ps, $\delta = 0.1$. The two-level system and structural parameters are are: $\xi \in [0,\ 108]$, $d = 0.2$, $\tau_p = 4.69$ fs, $\delta = 0.25$, $\Omega_0 = 1.49$ fs^{-1}.

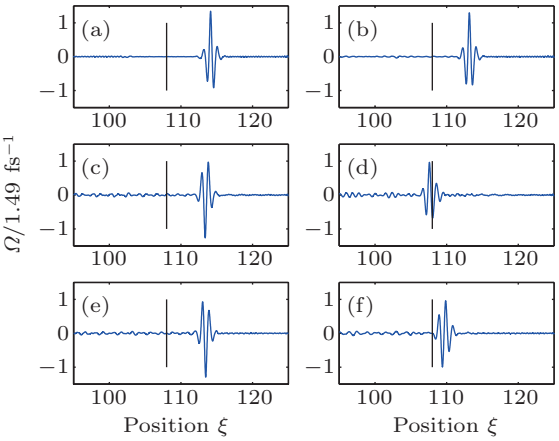


Fig. 4. (color online) Electric-field profiles for different chirp parameters at time $t = 0.376$ ps. (a) $\alpha = 0$ and $\delta = 0.1$; (b) $\alpha = 0$ and $\delta = 0.25$; (c) $\alpha = -0.05$ and $\delta = 0.1$; (d) $\alpha = -0.05$ and $\delta = 0.25$; (e) $\alpha = 0.05$ and $\delta = 0.1$; (f) $\alpha = 0.05$ and $\delta = 0.25$. The two-level system and structural parameters are: $\xi \in [0,\ 108]$, $d = 0.2$, $\tau_p = 4.69$ fs, $\delta = 0.25$, $\Omega_0 = 1.49$ fs^{-1}.

In the next, we will investigate the effect of the structure's length and the slab number N_s on the pulse shaping. In Fig. 5, the electric field profiles of the penetrating pulses are shown with different lengths and periods of the subwavelength structures. In our calculation, we assumed that one detector is located as $\xi = 50$ and record the electric field profiles of the penetrating pulses. The parameters is the same to in Fig. 4, but the density is $d = 0.2$. The left plots in Fig. 5 correspond

to the case of a compact structure ($\delta = 0.1$). Here we find that a single-cycle pulse can be gotten after propagating through a short structure. From the right plots in Fig. 5, we can conclude that the compact structure is more effective to split the few-cycle pulses because more energy of the initial pulse enters into the structure. The electric field profiles with a bigger amplitude and smaller duration in time coordinate are shown in Fig. 6 when the slab number is changing. In Fig. 6, we observe that the pulse with one optical cycle can be shaped into a shorter pulse with optical cycle number 0.3 after propagating through a subwavelength structure which consists of 30 slabs with 0.25 cycle. Based on the results above, a smaller shaper of the few-cycle pulse can be designed.

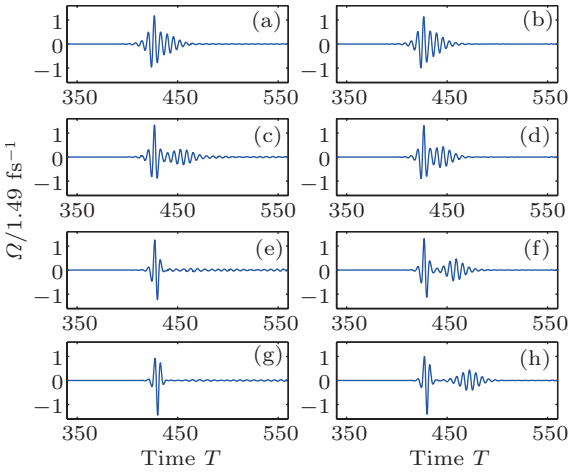


Fig. 5. (color online) Electric-field profiles in time coordinates for $\xi = 50$. As a comparison, the same structure lengths are considered: (a) $L = 1$, $\delta = 0.25$; (b) $L = 1$, $\delta = 0.1$; (c) $L = 2$, $\delta = 0.25$; (d) $L = 2$, $\delta = 0.1$; (e) $L = 4$, $\delta = 0.25$; (f) $L = 4$, $\delta = 0.1$; (g) $L = 6$, $\delta = 0.25$; (h) $L = 6$, $\delta = 0.1$. The other parameters are given here: $d = 0.2$, $\Omega_0 = 1.49$ fs^{-1}, $\delta = 0.25$, $\tau_p = 4.69$ fs. The initial cycle number N_{FWHM} and envelope area A are 1.83 and 4π, respectively.

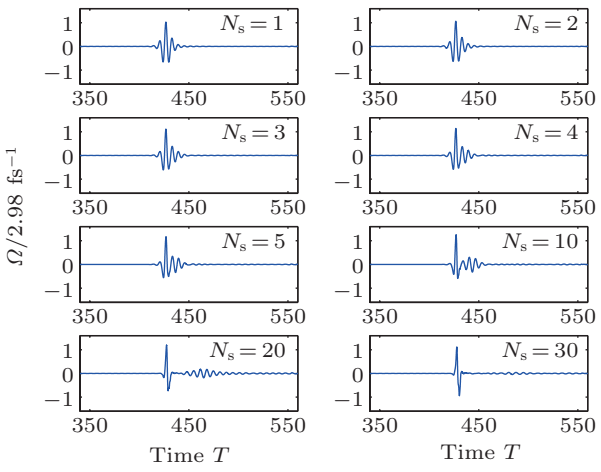


Fig. 6. (color online) Electric-field profiles in time coordinate for $\xi = 50$. Here, different slab numbers N_s are considered, which have been labeled in different subfigures. The other two-level system and structural parameters are given here: $d = 0.2$, $\Omega_0 = 2.98$ fs^{-1}, $\delta = 0.25$, $\tau_p = 2.35$ fs. The initial cycle number N_{FWHM} and envelope area A are 0.92 and 4π, respectively. The cycle number of the output field is about 0.3 when $N_s = 30$.

At last, we also calculate the pulse propagation of few-cycle pulses in a subwavelength structure with a non-ideal period. In these simulations, the lengths of slabs have a five percent error. After analyzing an amount of numerical calculations, we find that penetrating pulses with the same duration can be gotten. If the error is larger, the difference of the output fields between ideal and non-ideal periodic structures becomes obvious and can be observed.

4. Summary

In summary, the shaping of few-cycle pulse via a subwavelength structure is studied theoretically by using the full Maxwell–Bloch equations beyond the traditional frame of SVEA and RWA. The impacts of some parameters, such as media density, layer thickness, layer length, pulse duration, chirp coefficient, and so on, on the output fields were investigated. We found that the subwavelength structure with a compact period is more efficient to shape few-cycle pulses into shorter laser pulses. Such structure suggests a way toward shortening high-intensity laser fields, which might help us achieve ultrashort pulses with single-cycle or even subcycle duration.

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