Generation of relativistic high-order-mode laser pulse using plasma waveguide

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Keywords: laser-plasma interaction, high-order-mode laser pulse, mode conversion, plasma waveguide

Abstract
An all-optical method for generating ultra-intense high-order-mode light pulse is investigated with three-dimensional particle-in-cell simulation. We find that the conversion from a short intense circularly polarized incident Gaussian laser pulse into a transverse magnetic (TM) mode occurs as it propagates into a micro plasma waveguide. The strength of the longitudinal electric field of the excited TM modes can be almost two orders of magnitude higher than that of the original laser. The simulation results show that, for the lower-order modes, the trapped electrons lead to their revolving transverse structures. A linear plasma waveguide model is presented to predict the mode pattern and intensity of the longitudinal electric fields, which are in excellent agreement with those in the simulations. Relativistic-intense high-order-mode light can be useful for many applications, including accelerating charged particles to high energies.

1. Introduction
Existing petawatt table-top lasers can deliver femtosecond light pulses with intensity far beyond $10^{18}$ W m$^{-2}$ [1]. Their nonlinear interaction with plasma can create longitudinal electric fields with very strong gradients that are more than 10$^3$ times larger than that in the conventional accelerators [2, 3]. Such laser-driven accelerators can generate MeV, or even GeV, energetic charged particles within a centimeter or less [2, 3] and have thereby opened the door for exploring high energy density physics and astrophysics in the laboratory. Recently, it was shown that intense short wavelength optical modes with inherent longitudinal electric fields are particularly effective for electron acceleration [4–7]. The resulting electrons have typically superponderomotive temperature and structures that are suitable for generating bright x-rays [4], energetic ions [5], etc. However, realizing such high-order short wavelength light in the laboratory is still challenging, mainly because of radiation damage of the optical modulators.

Plasmas have a distinct advantage for manipulating intense laser light because they have no thermal damage threshold. Currently, several plasma-based methods have been proposed for generating intense laser light with strong longitudinal wave electric fields, including simulated Raman scattering (SRS) [8], hologram grating [9], spiral phase plates [10], etc. However, SRS functions only for rarefied plasmas and not-so-intense lasers, since the gratings and phase plates can be easily smoothed out by the prepulse or amplified spontaneous emission of the incident laser [8–10]. Another plasma-based method is to use hollow or plasma filled channels [11, 12] originally proposed for eliminating diffraction spreading of laser light [11]. The laser pulse can be guided by the channel for several tens Rayleigh lengths, but the laser intensity is limited to $10^{13}–10^{14}$ W cm$^{-2}$ [11].

In this paper, we propose an all-optical method for generating relativistic ($>$10$^{18}$ W cm$^{-2}$) high-order-mode optical waves using cylindrical micro plasma waveguide. Our 3D-PIC simulations show that a suitable hollow waveguide partially filled with electrons pulled out by the intense light of a circularly polarized (CP) Gaussian
laser pulse propagating in it, can act as a mode convertor for the latter into the transverse magnetic (TM) modes of the waveguide. The high-order TM modes have structured longitudinal electric fields suitable for accelerating charged particles. Because the group velocity of the modes decreases with the mode order, the order of the modes increases from the front to the tail of the pulse until the highest-order eigenmode of the plasma waveguide is reached, and their longitudinal electric field can be of the same order as that of the oscillating electric field of the intense incident laser. It is also found that electrons pulled out can be effectively trapped in the wave field of the lower-order modes and cause their rotation. A linear plasma-waveguide theory is established, which can be extended to predict the spatial distribution and intensity of the TM-mode longitudinal electric fields. The theoretical results are in excellent agreement with in the simulations.

2. Linear plasma waveguide model

We consider a relativistic Gaussian laser pulse propagating in a hollow cylindrical waveguide with overdense walls, as shown in figure 1(a). The normalized laser amplitude $a_0 = eE_0/m_\text{e}\omega_0\epsilon$, where $E_0$ is the peak laser electric field, $\omega_0$ is the central laser frequency, and $m_\text{e}$ and $-\epsilon$ are the electron rest mass and charge, is much larger than unity. The channel wall is opaque if $\omega_{pe} > \omega_0$ [13], where $\omega_{pe} = (4\pi\epsilon_0n_e/m_\text{e})^{1/2}$ is the relativistic plasma frequency and $n_e$ is the electron density in the waveguide wall. The relativistic factor can be expressed as $\gamma = [1 + a_0^2 s^2(r_0)]^{1/2}$, where $s(r_0) = \exp(-r_0^2/\sigma_0^2)$ is the radial profile of the light wave envelope, $r_0$ is the inner radius of the waveguide, and $\sigma_0$ is the spot radius of the laser. Thus, most of the laser energy will be reflected by the wall if

$$\frac{n_e}{n_r} \gg \sqrt{1 + a_0^2 \exp(-2r_0^2/\sigma_0^2)},$$

where $n_e = m_\text{e}\omega_0^2/4\pi\epsilon^2$ is the critical plasma density. In this case the plasma waveguide can be regarded as a traditional optical waveguide. The solution for the TM$_{nl}$ mode electric field in the cylindrical coordinates $(x, \phi, r)$ can be expressed as [14]

$$E_{x}^{TM}(x, \phi, r) = k_l^2 W_0,$$

$$E_{\phi}^{TM}(x, \phi, r) = ik_x W_0/r,$$

$$E_{r}^{TM}(x, \phi, r) = ik_x \frac{\partial W_0}{\partial r},$$

where

$$W_0 = D \frac{J_0(r_0k_l)}{k_l^2} \cos(n\phi) \exp(-k_x x)$$

is the electric Helmholtz potential, $D$ is a constant coefficient determined by the incident laser pulse, $J_0(r_0k_l)$ is the first-kind nth order Bessel function, $k_l = \nu_{nl}/r_0$ and $k_x = (k_0^2 - k_l^2)^{1/2}$ are the radial and longitudinal wave numbers, $\nu_{nl}$ is the nth root of the nth order Bessel function of the first kind, $k_0 = \omega_0/c = 2\pi/\lambda_0$ is the wave number of the laser and $\lambda_0$ is the laser wavelength.

Since $k_x^2 = k_l^2 - k_r^2 \geq 0$, we obtain the minimum waveguide radius $r_0^0 = \lambda_0\nu_{nl}/2\pi$. For $n = 1$, the root $\nu_{11}$ of the TM$_{11}$ mode and $r_0^0$ are shown in table 1, where $\nu_{nl}$ is the nth root of the nth order Bessel function. For example, we have $r_0^0 = 2.62 \lambda_0$ if the TM$_{11}$ mode is wished to be excited. The quasi-steady-state mode structures from equation (2) for different $l$ are shown in figures 1(b1)–(b5). For laser intensities in the range $10^{18}$–$10^{22}$ W cm$^{-2}$, the waveguide loses its excitation effect [15]. We can therefore assume that the excitation ratio of the peak radial electric field is $R_{nl} = E_r/E_0 \sim 1$, where $E_r$ is the peak radial electric field inside plasma waveguide. Since $E_r/E_r \approx E_{x}^{TM}/E_{r}^{TM}$ [4, 5], the magnitude of the longitudinal electric field excited in the plasma waveguide can approximately expressed by

$$E_x \approx E_0 \frac{k_lj_{l, \text{max}}}{k_xj'_{l, \text{max}}} \exp(-r_0^0/\sigma_0^2) = E_0 \frac{1.164}{(2\pi\nu_{nl}/\nu_{11}\lambda_0)^2} \exp(-r_0^2/\sigma_0^2),$$

where $j_{l, \text{max}}$ and $j'_{l, \text{max}}$ are the maximum values of the 1st order Bessel function of the first kind and its derivative with respect to $r$, respectively. The field distribution ratio $R_{xx} \approx E_x/E_r$ between the longitudinal and radial components inside the plasma waveguide and the excitation $R_{xx} = (E_{x0} + E_x)/E_{x0}$ of the longitudinal electric field can be estimated from equation (6), where $E_{x0}$ is the initial peak longitudinal electric field. It should be noted that equations (2)–(6) are valid only for $\sigma \gg r_0$. 

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**Note:** The text contains mathematical equations and physical descriptions that are typical for a scientific paper. The content is meant to be read and understood by readers familiar with the fields of plasma physics and wave propagation in waveguides.
3. Simulation and discussion

We use the 3D PIC simulation code EPOCH [16] to explore the nonlinear propagation of a relativistic CP Gaussian laser pulse in a plasma waveguide. The profile of the laser pulse is \( a = a_0 \sin(\pi t/\tau_0) \exp(-r^2/\sigma_0^2) \), where \( a_0 = 6.65, \tau_0 = 10\tau_0, \sigma_0 = 3\lambda_0, \lambda_0 = 0.8 \mu m \) and \( T_0 = 2.67 \) fs, respectively. The corresponding laser intensity, power, and total energy are \( 1.89 \times 10^{20} \) W cm \(^{-2}\), 34 TW, and 0.9 J, respectively. A waveguide of radius \( r_0 = 3\lambda_0 \), length \( L_0 = 30\lambda_0 \), and thickness \( \Delta r = 0.5\lambda_0 \) is located between \( x_1 = 0 \) and \( x_2 = 30\lambda_0 \). The singly ionized gold plasma wall of the waveguide has initial electron and ion densities \( n_e = n_{Au} = 33.5n_0 \). In fact, any material with large atomic mass can be used as long as the plasma density \( n_e \) after ionization satisfies the

![Figure 1](image-url)

**Figure 1.** (a) Schematic of the laser interaction in the waveguide. The inset at the top-left corner shows the normalized electron density profile \( n_e/n_0 \) inside the channel at \( t = 25T_0 \) obtained from the PIC simulation. At this moment, the center of the original light pulse (if no interaction took place) would roughly be at \( x = 15\lambda_0 \). (b) Spatial evolution of the transverse structure of the longitudinal electric field of the TM\(_1l\) modes along the axial direction from the theoretical model (b1)–(b5) and from the PIC simulation (b6)–(b10) at \( x_l = 20\lambda_0, 16\lambda_0, 12\lambda_0, 8\lambda_0, \) and \( 4\lambda_0 \) respectively. The black dots in the panels (b6)–(b10) represent typical trapped electrons. The colorbars in (b) are for the panels (b6)–(b10) showing the mode electric fields (normalized by \( w_m/c \)). The color in the panels (b1)–(b5) is relative, obtained by selecting the constant \( D \) in equation (5) such that the peak intensity roughly matches that from the simulation.

| \( l \) | \( \nu_{1l} \) | \( r^*_m \) |
|---|---|---|
| 1 | 3.83 | 0.61 |
| 2 | 7.02 | 1.12 |
| 3 | 10.2 | 1.62 |
| 4 | 13.3 | 2.12 |
| 5 | 16.5 | 2.62 |
| 6 | 19.6 | 3.12 |
| 7 | 22.8 | 3.62 |
| 8 | 25.9 | 4.12 |

Table 1. The root \( \nu_{1l} \) of the 1st order Bessel function and the required minimum inner radius of the plasma waveguide \( r^*_m \) for the excitation of the \( l \)th order TM mode.
Near the channel wall the electron density is higher and the laser amplitude is low, so that the constant of the plasma waveguide is 

\[ \frac{\omega}{v_c} = \frac{\sqrt{k_0^2 - \nu_0^2/\omega_0} - \nu_0^2/\omega_0}{\omega_0}, \]  

which shows that it decreases with increase of \( l \), so that the higher-order modes will lag behind the lower-order ones. The action of the plasma waveguide for laser light here is analogous to that for microwave and optical-fiber waveguides: the input EM waves are converted into propagating eigenmodes of the plasma waveguide according to the boundary conditions, and the mode structures also depend on the properties of the waveguide material. These results here suggest that even in the highly nonlinear relativistic regime, the plasma waveguide can be used to convert laser light into the TM modes.

The top left inset in figure 1(a) shows the profile of the wave-trapped electrons inside the plasma waveguide. The electrons have been pulled out from the waveguide wall by the intense CP light-wave and self-consistently trapped in them, resulting in rotation of the modes, as can be seen in figures 1(b6)–(b10). Since the different modes are separating from each other along x, the trapped-electron distribution appears to be spiral-like. From the microscopic point of view, the process involves the transfer of (orbital) angular momentum of the electrons to the spin angular momentum of photons, similar to that in producing vortex light by helically moving electrons [7–10, 18, 19]. The higher order modes have less trapped electrons. As a result, the profiles of the higher order modes are nearly the same as that from the linear model, as can be seen in figures 1(b1)–(b5).

Although the light in the core of the waveguide is of high intensity, its periphery, which interacts with the waveguide wall, is of much lower intensity. Thus, the density of the wall electrons pulled into the waveguide channel and trapped in the wave fields is much lower than that in the wall. For the same reason, the radial light pressure which tends to drive the electrons back into the waveguide wall is small [20]. Moreover, since the mass of the Au ions in the waveguide is high, the space-charge field created by the displaced electrons is not sufficient to displace the ions. As a result, the nonlinear light wave-electron interaction is not sufficient to change the effective profile of the original waveguide wall, which is consistent with the observation that the simulated results are rather similar to that from the linear theory.

Local excitation of the incident light inside the plasma waveguide can be seen in figure 2 for the envelopes of the maximum radial and longitudinal laser and TM mode electric fields in the transverse planes along the axial direction, as well as the spectrum of the longitudinal electric field in the \( x, y \) plane at \( t = 30 T_0 \). The increase in \( E_z \) can be attributed to nonlinear focusing of light in the hollow waveguide channel. The corresponding dielectric constant of the plasma waveguide is [19]

\[ \varepsilon = 1 - \frac{n_e}{\sqrt{1 + a_0^2 n_e}}. \]  

Near the channel wall the electron density is higher and the laser amplitude is low, so that \( \varepsilon < 0 \) and the laser will be focused to the waveguide core, where \( \varepsilon > 0 \) [22, 23]. For a Gaussian laser pulse, the initial amplitude of the normalized longitudinal wave electric field \( E_z \) is 0.25, only about 1/30 of the radial wave field, as can be seen in figure 2(b). In the core of the waveguide, the peak value of \( E_z \) can be as high as 1.5, or 6 times, and the intensity 36 times that of the incident laser. That is, the magnitude of the converted longitudinal wave electric field of the TM mode can approach \( 1/4 \) that of the (mainly transverse wave field of the) incident laser. In the process, the spread of the longitudinal electric field is also broadened. This is because, as discussed earlier, high-order TM modes with smaller group velocities are excited in addition to the TM11 mode. Both the excitation and broadening of the longitudinal wave electric field are beneficial for axial acceleration of charged particles. One can also see in the \( E_z \) spectrum periodic small peaks at \( k_0 t \) intervals. These can be attributed to the much weaker harmonics of the intense incident laser light excited in the nonlinear laser-electron interaction when an ultra-intense laser pulse propagates into the plasma channel [24–27].

Which modes can be excited in the plasma waveguide depends on the radius \( r_0 \) of the plasma waveguide. If \( r_0 \ll \lambda_0/2 \), the waveguide entrance can be regarded as a pin-hole, and according to linear optics, light diffraction
should occur [28]. However, for relativistic lasers the pulled-out electron the radial extent of a laser wavelength can partially block the entrance, leading to an irregular spot distribution in the region I in figures 3(a)–(d) as well as discrepancy between the theoretical and simulation results. However, as \( r_0 \) approaches \( \sigma_0 \), the blocking effect is gradually weakened and the TM mode begins to dominate in the region II. As a result, high-order modes appear, as can be seen in figures 3(e)–(h). Figure 4(a) clearly shows excitation of the longitudinal wave electric field in this high-order-mode excitation region. Figures 3(h)–(n) shows that for still larger \( r_0 \) values mode excitation rapidly becomes weakened, thereby entering region III of no high-order-mode excitation. Note that in the regions II and III, where the waveguide mode dominates, \( I_{\text{max}}, R_{\sigma r}, R_{xx}, \) and \( R_{yy} \) are all in good agreement with that of the theoretical results in figure 4(a), as well as table 1. That is, the linear theory given by equations (2)–(6) can be used for estimating the properties of the nonlinear relativistic TM plasma waveguide modes.

Analysis of our PIC results also show that for \( r_0 = 3 \sigma_0 \), only 10% of original laser energy is lost in the nonlinear interaction, mostly to the wall electrons that are pulled out, trapped, and accelerated by light waves. Most of the laser energy is preserved in the radial (80%) and longitudinal (10%) components of the laser fields. With decrease of \( r_0 \), the laser-to-electron energy conversion efficiency \( \eta \) increases because of more effective electron acceleration by the enhanced \( E_0 \) field, as can seen in equation (6). For given \( r_0 \) and \( \sigma_0 \), \( \eta \) is determined by \( a_0 \). Both the radial electron momentum \( p_r \) and the longitudinal electron momentum \( p_z \) from the high-order-mode electric field \( E_z \) are proportional to \( a_0 \) [5]. For the contribution of the laser ponderomotive force to \( p_z \), the superluminous phase velocity of the wave in the plasma waveguide must be taken into account if \( E_0 > a_x \), where \( a_x = [2\pi c/(\nu_{ph} - c)]^{1/2} \) is the critical amplitude and \( \nu_{ph} \) is the light phase velocity [5]. In fact, it changes from \( a_x^2 m_e c/2 \) for \( a_0 < a_x \) to \( a_0^2 m_e c/(\nu_{ph}^2 - 1)^{1/2} \) as \( a_0 \) is further increased. In addition, we found that the density of the pulled-out wall electrons scales as \( a_0^4 \), where \( \chi \) is related to the ionization degree of the wall material [5]. Therefore \( a_0 \) is closely related to \( \eta \) and therefore can affect the energy conversion efficiency from the laser to the TM modes.

Figure 4(b) shows that \( R_{yy} \) contains small fluctuations around 0.175, in agreement with equation (6), suggesting that the energy distribution for both field components inside a plasma waveguide of the given radius is fixed. Furthermore, \( E_z \) increases linearly with \( a_0 \), so that \( R_{yy} \approx 1 \) and there is not much excitation of the radial electric field inside the plasma waveguide for the laser intensity range \( I_0 = 10^{18} - 10^{22} \) W cm\(^{-2} \). For \( E_z = 0.33 E_0 \), the amplitude is of the same order as that of the radial one, and can be much higher than that of the Gaussian pulse. In figure 4(b), \( a_0 \) is limited to 60 because the theoretical model is only valid for that laser-intensity range. Radiation reaction effects will have to be taken into account if \( I_0 > 10^{22} \) W cm\(^{-2} \).

4. Other effects

So far we have considered mode conversion of a CP laser pulse. Figures 5(a) and (b) show the TM mode structures excited by a \( p \)-polarized laser pulse. One can see that they do not rotate, as expected, since the electrons oscillate only along the light polarization (here the \( y \) direction). However, one can see in table 2 that theendas

**Figure 2.** Envelope of the maximum (a) radial and (b) longitudinal wave electric fields in the transverse \((y, z)\) plane along the axial direction \( x \) at \( t = 30\tau_0 \). (c) Spectrum of the longitudinal electric field along \( r = 1\lambda_0 \) in the \((x, y)\) plane at \( t = 30\tau_0 \). The black (‘inc’) and red (‘pw’) curves are for the input Gaussian laser pulse and the plasma waveguide modes, respectively.
Figure 3. The highest-order mode structures of the longitudinal electric field excited in the plasma waveguide at different waveguide radius $r_0$. With increasing $r_0$, three regions appear: (a)–(d) the hole-blocking region I ($r_0 \ll \sigma_0$), (e)–(h) the high-order-mode excitation region II ($r_0 \sim \sigma_0$), and (i)–(l) the high-order-mode non-excitation region III ($r_0 \gg \sigma_0$), for given laser spot radius $\sigma_0 = 3\lambda_0$.

Figure 4. Dependence of the ratio $R_{xx}$ of the longitudinal-to-radial wave electric fields in the plasma waveguide on the (a) waveguide radius $r_0$ and (b) laser amplitude $a_0$. The blue curves and symbols in (a) are for $R_{xx}$ and $R_{rr}$ between the longitudinal and radial mode electric fields, respectively, inside the plasma waveguide and that of the Gaussian incident laser. The properties of the three regions (separated by vertical dashed lines) are in good agreement with that in figure 3. In (b), the longitudinal and radial wave electric fields in the plasma waveguide as functions of the laser amplitude $a_0$ from the PIC simulation and theory are given by the blue curves and symbols, respectively.
field strengths in the plasma waveguide are comparable with that excited by the CP laser. That is, relativistic LP laser pulses can also generate high-order waveguide modes.

It is well known that a conical waveguide can focus laser light [29, 30]. Table 2 for a conical waveguide with 5° half-angle and entrance radius 3λ0 shows that the peak magnitude of both the transverse and longitudinal electric field components in the waveguide are indeed higher than that of the cylindrical waveguide. In fact, the longitudinal electric field intensity can be 100 times that of the incident laser. The effect on the pulled-out electrons can also be clearly seen. However, l_{max} is less than that of a cylindrical waveguide because of the narrowing of the channel, and only up to the 4th order TM mode can appear, as can be seen in figures 5(c) and (d).

In practice, intense laser pulses are usually operated at non-zero incidence angle in order to avoid self-damage by the reflected light. We have thus also carried out simulations with a small incidence angle. Figures 5(e) and (f) and table 2 show that the mode conversion process investigated above still takes place, but the center of the high-order-mode structure can deviate slightly from the axial direction.

5. Experimental considerations and potential applications

The proposed scheme should be realizable in the laboratory with existing technology. First, ultrashort laser pulses with intensity up to 10^{22} W cm^{-2} and contrast of 10^{-10} are available [31–35]. Several techniques for
fabrication of submicron scale target are now available. For example, electrodeposition of metallic ions now allows a tailored target to grow to the required size inside porous template pores. In fact, a nanowire array of length 10 \( \mu \)m, width 2 \( \mu \)m, and aperture 400 nm has been successfully used for generating extreme pressures \[34, 35\]. Modern precision electrochemical etching can now also operate at the micron level \[36\]. Such and other novel methods can be useful for fabricating the hollow waveguides needed in our scheme. In fact, the use of a micro-channel target of diameter 5 \( \mu \)m has been successfully used in an electron acceleration experiment \[37\]. On the other hand, the ubiquitous laser-spot quivering in real situations may post a problem for the target aiming since the waveguide aperture here is comparable with the laser spot size. In this case one can use an appropriate slab target with an array of closely packed cylindrical holes serving as waveguides.

6. Potential applications

Plasma waveguide are suitable for modulation of \(10^{18}–10^{20}\) W cm\(^{-2}\) lasers. Tunability of the excited optical modes and intensities of their longitudinal electric fields make them suitable for accelerating externally injected or self-generated negatively charged particles. Existing work \[4\] found that, with a \(10^{20}\) W cm\(^{-2}\), 15 fs, 1.8 J laser, the electrons extracted from the walls can have a total charge of 2.4 nC and are accelerated to a maximum energy of about 300 MeV in a 400 \( \mu \)m long interaction. The asymmetric transverse fields of the high-order TM modes can also act as efficient wiggler fields for these energetic electrons, so that well-collimated bright emission of 1–100 keV hard x-rays are generated. Plasma waveguides can thus serve as compact bright x-ray sources. With higher laser intensity, a large amount of high-energy gamma photons can be generated via Compton backscattering. If collided with a counter-propagating laser pulse, the accelerated electrons may lead to pair production via the Breit–Wheeler process \[38\]. Furthermore, focused energetic electrons, useful for target normal sheath acceleration, can be achieved by using a proper channel size \[3, 39\]. It is also of interest to note that there is no change of the central frequency of the incident laser as it is converted into the lowest-order waveguide mode. Accordingly, high-order terahertz TM modes, useful for high-gradient terahertz-driven linear electron accelerators \[40\], should be realizable if the injected laser is of terahertz frequency.

7. Conclusion

In summary, we have proposed the use of miniature plasma waveguide as optical convertor for ultra-intense Gaussian laser light. Our 3D simulation results show that the excited waveguide TM modes can be of high-order, and difference among the mode group velocities lead to their separation along the waveguide, with the faster low-order modes in the front. In the high-order mode excitation region, the intensity of the local longitudinal mode electric field can be one order of magnitude higher than the initial value. Our simulation results agree quite well with that of a linear theoretical model, suggesting that the latter can be used for estimating the effects of the plasma waveguide. It is also shown that the optical mode conversion by plasma waveguide considered here is quite robust and it tolerates different laser intensities, laser polarizations, as well as (small) angles of incidence. The required laser intensities and precision waveguides needed here are also attainable with available petawatt lasers and target technologies, thereby offering a realizable approach for accelerating charged particles to high energies and generating bright x-rays.

Acknowledgments

This work is supported by the National Key R&D Program of China (Grants No. 2016YFA0401100 and 2018YFA0404802), the National Natural Science Foundation of China (Grants No. 11705280, No. 11774430 and No. 11622547), the Research Project of NUDT (Grant Nos. ZK18-03-09 and ZK18-02-02), the Science Challenge Project (Grant No. TZ2016005), the National Basic Research Program of China (Grant No. 2013CBA01504) and the State Key Laboratory of High Field Laser Physics at SIOM.

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