Electromagnetically induced guiding of counterpropagating lasers in plasmas

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The interaction of counterpropagating laser pulses in a plasma is considered. When the frequencies of the two lasers are close, nonlinear modification of the refraction index results in the mutual focusing of the two beams. A short (of order of the plasma period) laser pulse can also be nonlinearly focused by a long counterpropagating beam which extends over the entire guiding length. This phenomenon of electromagnetically induced guiding can be utilized in laser-driven plasma accelerators. [S1063-651X(99)00201-9]

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I. INTRODUCTION

The past decade witnessed a dramatic increase in intensities of laboratory-scale laser systems [1,2], stimulating significant interest in the nonlinear optics of plasmas [3–5]. One of the areas of active research is propagation of intense laser pulses over long distances in plasmas because of its importance for a number of applications, including laser wakefield particle accelerators [6]. A laser will, in free space, remain focused over a diffraction length (Rayleigh range) $Z_r = \pi\sigma^2/\lambda_0$, where λ_0 is the laser wavelength, and σ is the laser spot size at the focus. Nonlinear effects, such as relativistic and ponderomotive self-focusing, which can overcome diffraction, have been studied theoretically [7,8] and observed experimentally [9]. Nonlinear self-focusing is ineffective for short laser pulses [10] and requires a very high laser power $P_{c1} = 17\omega_0^2/\omega_p^2$ GW, where $\omega_p^2 = 4\pi n_0 e^2/m$ is the plasma frequency, n_0 is the background electron density, and $-e$ and m are the electron density and mass, respectively.

In this paper we demonstrate how a laser pulse can be guided through the plasma without diffraction due to its nonlinear interaction with another, counterpropagating (guiding) beam. Such electromagnetically induced guiding (EIG) occurs at laser intensities much below the relativistic threshold P_{c1}. The guiding and guided beams interfere, forming a periodic intensity pattern which ponderomotively drives the plasma wave. The guiding beam undergoes a Bragg reflection off this periodic density perturbation, scattering into the guided pulse. Two distinct problems are analyzed in this paper: (i) mutual guiding of two long counterpropagating laser pulses and (ii) guiding of an ultrashort tightly focused laser pulse by a counterpropagating, lower-intensity Bessel beam. In both cases the length of the long (guiding) beam is approximately twice the desired propagation distance.

The technique of employing a second laser pulse to change the propagation properties of the first pulse is widely known in conventional (atomic) nonlinear optics. For example, by utilizing the effect of electromagnetically induced transparency (EIT) [11], the medium which is opaque to the laser pulse *taken alone* can become transparent in the presence of the second pulse. As was recently demonstrated by Harris [4], a similar process can take place in cold electron plasmas, where the two lasers have to be detuned by $\Delta\omega \approx \omega_p$. The transverse profile of the nonlinear index of refraction can be changed by the second pulse to prevent, e.g., an intense pulse from self-filamenting [12]. EIG of *copropagating* lasers in plasma was also considered by Gibbon and Bell [13] in the context of the cascade focusing, and by other authors [14]; the required laser power in that case is, however, much larger than in the counterpropagating case, and guiding of pulses shorter than a plasma wavelength is impossible.

The remainder of the paper is organized as follows. Section II introduces the general formalism describing the interaction between two low-intensity counterpropagating lasers. The results of a direct numerical simulation, confirming the mutual focusing of two long pulses, are also presented. In Sec. III we develop a simplified theory of EIG by assuming that the guided pulse is transversely Gaussian. Section IV deals with guiding ultrashort pulses. It turns out that the most power-efficient way of guiding a short (of order plasma period) pulse is to employ a long (twice the desired propagation distance) low-intensity Bessel beam. It is demonstrated that, depending on the frequency detuning of the two lasers, either the zeroth- or the first-order Bessel beam is appropriate for guiding. The validity of the linear theory is also discussed. Section V concludes and outlines the directions for future work.

II. BASIC FORMALISM

Consider two circularly polarized counterpropagating laser pulses $\vec{a}_0$ and $\vec{a}_1$, where $\vec{a}_{0,1} = a_{0,1}/2(\vec{e}_x \pm i\vec{e}_y)e^{i\theta_{0,1}} + \text{c.c.}$ are the normalized vector potentials: $a_{0,1}(\vec{r}_\perp, z, t) = eA_{0,1}/mc^2$. The subscripts 0 and 1 distinguish the right-moving (guided) pulse and the left-moving (guiding) pulses, respectively. The phases of the waves are $\theta_0 = (k_0 z - \omega_0 t)$ and $\theta_1 = (k_1 z + \omega_1 t)$, and we choose $|\Delta\omega| = |\omega_0 - \omega_1| \ll \omega_0$. It is further assumed that the pulses are propagating through a tenuous plasma $\omega_p \ll \omega_{0,1}$, so that the phase and group velocities of both pulses are close to the speed of light c, and, to lowest order in $\omega_p/\omega_0, k_0 \approx k_1 \approx \omega_0/c$.

Assuming that both lasers are of nonrelativistic intensity, $a_{0,1} \ll 1$, the dominant nonlinear force experienced by the plasma electrons is the ponderomotive $\vec{v}_1 \times \vec{B}_0 + \vec{v}_0 \times \vec{B}_1$ force, where $\vec{v}_{0,1} = c\vec{a}_{0,1}$. On a time scale much longer than $1/\omega_0$, plasma electrons are pushed by the intensity gradient of the "optical lattice" which is produced by the interference of the two beams and has a spatial periodicity $k_0 + k_1 \approx 2k_0$. This ponderomotive force generates a density perturbation $\delta n/n_0 = \hat{n}\ \exp i(\theta_0 + \theta_1) + \text{c.c.}$, which couples the two laser beams. Serving as an index grating with the wave number $k = k_0 + k_1$, oscillating at the difference frequency $\Delta\omega$, this density perturbation backscatters the left-going pulse a_1 into the right-going pulse a_0, and *vice versa*. This is precisely the mechanism responsible for the stimulated Raman backscattering in the plasma, which received a lot of theoretical [15,16] and experimental [17,18] attention. Below we demonstrate how this interaction between the counterpropagating lasers leads to the nonlinear focusing of one (or both) pulses.

The nonlinear interaction between a_0 and a_1 is calculated by substituting the modified plasma density into the paraxial wave equations for the right- and left-moving pulses. For example, for the guided pulse we obtain

$$2ik_0\left(\frac{\partial}{\partial z} + \frac{1}{c}\frac{\partial}{\partial t}\right)a_0 + \nabla_\perp^2 a_0 = -k_p^2 \chi_0 a_0, \tag{1}$$

where χ_0 is the mutual refraction index (witnessed by the guided beam), which is given by

$$\chi_0(\tau) = \frac{2\omega_0^2 |a_1|^2}{\omega_p} \int_0^\infty du\ \sin(\omega_p u) e^{i\Delta\omega u} \frac{a_0(\tau - u)}{a_0(\tau)}, \tag{2}$$

where $\tau = t - z/c$ is comoving with the guided pulse coordinate and the guiding pulse is assumed very long.

Equation (2) is obtained by formally solving the linear equation for the ponderomotive excitation of the density wake [19]:

$$\left(\frac{\partial^2}{\partial\tau^2} + \omega_p^2\right)\frac{\delta n}{n_0} = \frac{c^2}{2}\nabla^2 |\vec{a}|^2/2, \tag{3}$$

where $|\vec{a}|^2 = |\vec{a}_0|^2 + |\vec{a}_1|^2 + 2\vec{a}_1 \cdot \vec{a}_0$ is the total laser intensity and $\nabla^2 \approx (k_0 + k_1)^2$ for the counterpropagating lasers. Therefore, the magnitude of the density perturbation scales, roughly, as $\delta n/n_0 \sim a_0 a_1 \omega_0^2/\omega_p^2$. Therefore, Eq. (3) makes the advantage of using counterpropagating pulses apparent: the longitudinal gradient scale of laser intensity becomes very small (about $\lambda_0/2$), greatly reducing the laser power required to produce a given density modulation. Equation (3) is also a convenient starting point for comparing the mutual focusing of counterpropagating beams, considered here, and the earlier suggested [13,14] cascade (or beatwave) focusing of two *copropagating* beams. In the case of copropagation of two lasers with very close frequencies, one should use $c^2\nabla^2 \approx \partial^2/\partial\tau^2$. If the laser detuning is not close to the plasma frequency, $|\Delta\omega^2 - \omega_p^2| \sim \omega_p^2$, density perturbation scales as $\delta n/n_0 \sim a_0 a_1$. Relativistic corrections to the electron mass, which enters the expression for the plasma frequency, scales the same way. Plasma response can be increased, of course, by tuning the lasers to the plasma resonance [13], $|\Delta\omega| \approx \omega_p$, but this imposes additional restrictions on the plasma inhomogeneity. Comparing the two expressions for the density modulation demonstrates that counterpropagation enables a larger modification of the refraction index than copropagation.

As Eq. (2) indicates, the EIG of a given longitudinal laser slice τ is determined by the laser intensity at all earlier instances $\tau' < \tau$, so that the degree of focusing experienced by a short pulse is not uniform along the pulse and determined by its longitudinal profile. One finds that, if the pulse intensity drops off faster than exponentially in $|\tau|$, the ratio $a_0(\tau - u)/a_0(\tau)$ decreases with $|\tau|$. For example, for a longitudinally Gaussian profile $a_0 = \alpha_0 \exp(-\tau^2/2\tau_L^2)$ the leading edge $\tau < -\tau_L$ slowly erodes because $\chi_0(\tau) \propto \tau_L^2/\tau^2$. However, for $a_0 = \alpha_0 \text{sec}(\tau/\tau_L)$ (exponential decay in $|\tau|$) the leading edge is uniformly focused: for $\Delta\omega = 0$,

$$\chi_0 = 2|a_1|^2 \omega_0^2 \tau_L^2/(1 + \omega_p^2 \tau_L^2). \tag{4}$$

Since the plasma does not respond on a time scale faster than $1/\omega_p, \chi_0 \sim \tau_L^2$, resulting in the reduced guiding for ultrashort pulses. However, the leading edge remains exponential for all times.

The mutual refraction index χ_0 is, in general, a complex number. The real part of χ_0 is much larger than the imaginary part when the beams are detuned far from the plasma resonance, $|\Delta\omega^2 - \omega_p^2| \gg \omega_p^2$. For such detunings there is no energy exchange between beams, and their interaction is purely refractive. This is assumed for the rest of this section. The nonlinear interaction between a_0 and a_1 can then lead to electromagnetically induced guiding (of one or both beams) if χ_0 is peaked on axis. The refraction index depends on the spatial and temporal profiles of the pulses and their frequency detuning. For a long guided pulse of duration $\tau_L \gg \min(1/\Delta\omega, 1/\omega_p)$, the nonlinear index of refraction is independent of τ, ensuring uniform focusing of the entire guided pulse: $\chi_0 = 2|a_1|^2 \omega_0^2/(\omega_p^2 - \Delta\omega^2)$. Thus, the mutual refraction index acquires a transverse variation proportional to the transverse profile of the guiding laser. Two long transversely Gaussian beams focus each other if $\Delta\omega < \omega_p$.

III. NUMERICAL CONFIRMATION OF EIG AND THE ENVELOPE EQUATION

Since this type of nonlinear guiding relies on a rather delicate mechanism of generating an index grating with a very short wavelength, and then backscattering off this grating, we use direct particle in cell (PIC) simulation in a slab geometry to demonstrate the existence of the mutual focusing effect. Then we develop a simplified analytical description of the nonlinear guiding of Gaussian pulses in the slab and cylindrical geometries, interpret the simulation results, and address problem (ii).

To simulate the mutual focusing of two identical laser pulses, we use a 2D version of the relativistic electromagnetic PIC code VLPL (Virtual Laser Plasma Lab) [20] running on a single processor workstation. The grid size was 4000×120 with 4×10^6 electrons on it. The two counterpropagating laser pulses with wavelengths $\lambda_1 = \lambda_0 = 1\ \mu\text{m}$ are focused to the spot sizes $\sigma_0 = \sigma_1 = 3\ \mu\text{m}$ at their corre-

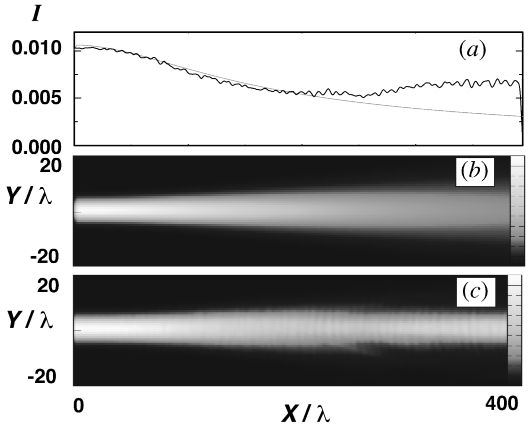

FIG. 1. (a) On-axis intensity $I=2|a|^2$ of a right-going laser focused at $z=0$ in vacuum (dashed line) and in the plasma (solid line). Contour plot of the intensity of the right-going laser in vacuum (b) and in the plasma (c).

sponding entrances into a 400 μm$\times$40 μm slab of $n_0 = 10^{19}$ cm^{-3} plasma. The normalized vector potentials of both lasers at their respective foci are equal to $\alpha_0=0.07$. The intensity gray scale plots for the right-moving pulse *in vacuo* and in the plasma are shown in Figs. 1(b) and 1(c), respectively. The snapshots at $t=500\lambda_0/c$ of the on-axis $(x=0)$ laser intensities for both cases are plotted versus the propagation distance in Fig. 1(a).

As Fig. 1(a) demonstrates, in the presence of the plasma the intensity of the right-going laser $|a_0|^2(z=L_g)$ is increased by a factor 2. Thus, the nonlinear interaction between the lasers strongly reduces laser spreading, confirming the electromagnetically induced guiding. Note that in this example EIG occurs at a very modest laser power, which is almost three orders of magnitude smaller than the power threshold for the relativistic focusing: $P/P_{c1}=0.0016$. The simulation was run until $t=800\lambda_0/c$, with no degradation of the mutual focusing. The driven density modulation remains coherent despite the modest plasma heating, which raises the temperature to about 500 eV. Note that for the parameters of the simulation, the growth rate of the Raman backscattering (RBS) instability of the guiding beam $\gamma_{\mathrm{RBS}} \ll \omega_p$, so that the frequency of the guided beam is outside of the RBS amplification bandwidth. RBS can grow from plasma noise, causing the beam reflection from the plasma. However, the saturated reflectivity is quite small for the underdense plasma assumed in the simulation [16].

A simplified analytical treatment of electromagnetically induced guiding can be developed by assuming that the guided laser pulse has a Gaussian transverse profile when it enters the plasma, and that it remains such in the plasma. Flat and cylindrical beams are considered separately because of the different laser intensity normalizations. The case of flat beams is important because it provides the basis for comparison with the results of the 2D PIC simulation. The guided beam is assumed to have an intensity profile given by $|a|_0^2 = (1/R_0)\alpha_0^2 \exp[-x^2/2R_0^2\sigma^2]$ in slab geometry and $|a|_0^2 = (1/R_0^2)\alpha_0^2 \exp[-r^2/R_0^2\sigma^2]$ in cylindrical geometry, where R_0 is the dimensionless spot size. Applying the source-dependent expansion to the envelope equations for a_0 similarly to the way it was done in Ref. [8], the equation for the normalized radius of a guided beam is derived:

$$\frac{d^2R_0}{d\overline{z}^2} = \frac{1}{R_0^3} + \frac{2k_p^2\sigma^2}{\sqrt{2\pi}R_0^2}\int_{-\infty}^{+\infty} dy\, y e^{-y^2/2R_0^2}\frac{\partial\chi_0}{\partial y} \qquad (5)$$

for flat beams, and

$$\frac{d^2R_0}{d\overline{z}^2} = \frac{4}{R_0^3} + \frac{4k_p^2\sigma^2}{R_0^3}\int_0^{+\infty} d\rho\, \rho^2 e^{-\rho^2/R_0^2}\frac{\partial\chi_0}{\partial\rho} \qquad (6)$$

for round beams, where $y=x/\sigma$ and $\rho=r/\sigma$ are the normalized transverse coordinates and $d/d\overline{z}=2k_0\sigma^2(\partial/\partial z + \partial/c\partial t)$.

To interpret the results of the PIC simulation, assume that the guiding pulse is also Gaussian with radius R_1, simplifying Eq. (5) to yield

$$\frac{d^2R_0}{d\overline{z}^2} = \frac{1}{R_0^3} - \frac{4k_p^2\sigma^2\alpha_0^2\omega_0^2}{\omega_p^2-\Delta\omega^2}\frac{R_0}{(R_0^2+R_1^2)^{3/2}}, \qquad (7)$$

confirming that the lasers focus each other if $\Delta\omega<\omega_p$. According to Eq. (7), the two lasers can, in principle, form a mutually guided state $R_0=R_1=1$ if $\alpha_0\approx 0.05$. This state is, however, unstable [21], so that it is never reached in a time-dependent situation, such as the one modeled by the PIC simulation. At the beginning of the plasma region the guided pulse interacts with the diffracted guiding pulse of much lower intensity. This explains why a somewhat higher $\alpha_0 = 0.07$ was needed in the simulation to observe the EIG.

IV. FOCUSING OF ULTRASHORT PULSES

We now address the second problem: focusing of an ultrashort tightly focused laser pulse by a counterpropagating lower intensity long laser pulse. One application of such ultrashort pulses is laser-wakefield acceleration. To increase the final energy of accelerated electrons, two quantities have to be maximized: the amplitude of the plasma wave left behind the pulse and the total acceleration length. Large plasma wake is excited by a laser pulse of a duration shorter than the plasma period. Assuming that the guided pulse has a longitudinal profile $a_0=\alpha_0 \mathrm{sec}(\tau/\tau_L)$, it can be demonstrated that the optimal pulse duration for wake excitation is $\tau_L = 1.2/\omega_p$. Guiding of such short pulses can be achieved by generating a plasma density depression on axis through the hydrodynamic expansion of laser-produced plasmas [22] or in a capillary discharge [23]. Here we demonstrate how an ''electromagnetic channel'' can be created by a counterpropagating laser beam.

The ultrashort pulse is guided by the long guiding pulse along its entire passage through the plasma. In contrast, the guiding pulse is only briefly affected by the guided pulse. Therefore, it has to propagate through the entire plasma region without any additional nonlinear focusing by the short pulse. We found that the most power-efficient choice of the guiding beam is a Bessel beam.

It has been known for some time [24] that apertured Bessel beams with sharply peaked radial profiles propagate without diffraction over distances much exceeding the Rayleigh length. Such beams transport energy within a narrow spot $\sigma_1 \ll W$ over the distance of order $L_g = 2\pi W \sigma_1/\lambda_1$, where the laser intensity profile is given by $|a_1|^2 = \alpha_1^2 J_p^2(r/\sigma_1)$ and the beam is apertured at the radius $W \gg \sigma_1$ [24]. The total power of a Bessel beam is given by $P_{Bp}/P_0 = \pi W \sigma_1 |\alpha_1|^2/\lambda_0^2$, where $P_0 = mc^3/r_e = 8.0$ GW and $r_e = e^2/mc^2$ is the classical electron radius. Therefore, the propagation distance is proportional to the beam power. Below we demonstrate how low-intensity Bessel beams can be used to guide ultraintense short Gaussian beams. Apart from the basic interest to nonlinear laser-plasma science, an experimental implementation of this novel guiding scheme would enable, for example, detailed studies of the stability of guided laser propagation in overmoded channels [25].

A large spot-size Gaussian beam, such that its Rayleigh length is equal to the required guiding distance L_g, may appear to be an alternative to the Bessel beam. At first glance, the on-axis mutual refraction index χ_0, which is proportional to the peak intensity of the guiding beam, would be, roughly, the same for the equal power Gaussian and Bessel beams. This is because the product of the peak on-axis intensity and the propagation distance is, approximately, the same for the equal power Gaussian and Bessel beams [26]. However, following Eq. (6), the focusing strength of the guiding beam is proportional to the *curvature* of χ_0, and is, therefore, larger for a Bessel beam.

The choice of p, the order of the Bessel beam, depends on the frequency detuning of the lasers. Consistent with our assumption of purely refractive interaction between beams, we consider two cases: (i) zero detuning $\Delta\omega = 0, p = 0$ and (ii) $\Delta\omega \gg \omega_p, p = 1$. For a zero-detuning case, assuming $a_0 = \alpha_0 \sec(\tau/\tau_L)$, χ_0 at the leading edge is given by Eq. (4), $\chi_0 = 2|a_1|^2 \omega_0^2 \tau_L^2/(1+\omega_p^2\tau_L^2)$. Clearly, χ_0 peaks on axis for a zeroth-order Bessel beam.

Numerical modeling of guiding a short pulse by the Bessel beam requires a 3D PIC simulation and imposes very strict limitations on the spatial and temporal resolution and the size of the simulation domain. This makes direct PIC simulations very difficult, and we rely on our analytical results. The power of the backward-moving Bessel beam P_{B0}, required for guiding, is estimated by substituting $|a_1|^2 = \alpha_1^2 J_0^2(r/\sigma_1)$ into Eq. (6) and assuming $R_0 = 1$:

$$k_0^2 \sigma^2 |\alpha_1|^2 \frac{2\omega_p^2 \tau_L^2}{1+\omega_p^2\tau_L^2} \left[t \frac{d}{dt} [e^{-t} I_0(t)] \right] = -1, \tag{8}$$

where $t = \sigma^2/2\sigma_1^2$ and $I_0(t)$ is a modified Bessel function. The strongest focusing occurs for $\sigma_1 \approx 0.8\sigma$. Since the guiding distance L_g is related to the beam power P_{B0}, Eq. (8) can be rewritten as a power threshold for the EIG:

$$P_{B0} = 1.27 \ \ \mathrm{GW} \frac{1+\omega_p^2\tau_L^2}{\omega_p^2\tau_L^2} \frac{L_g}{k_0\sigma^2}. \tag{9}$$

The EIG threshold P_{B0} differs from the relativistic guiding threshold P_{c1} in two respects: first, it is *independent* of the plasma density, enabling electromagnetically induced channeling in very tenuous plasmas, and second, it depends on the propagation distance. The qualitative reason for the first difference is that, although smaller plasma density reduces the effectiveness of the nonlinear guiding for the same fractional density perturbation $\hat{n}$, a smaller laser intensity is needed to generate the same $\hat{n}$, according to Eq. (3). The theory of the EIG, developed in this paper, assumes a linear plasma response, which is only valid when $\hat{n} < 1$. Combining this restriction with the guiding condition, given by Eq. (8), imposes an upper limit on the intensity of the guided pulse: $\alpha_0 \leq \lambda_0 \sigma/\lambda_p^2$. More extensive numerical study is needed to fully understand how the guiding is affected when the linear theory breaks down, but it is reasonable to expect that the guiding saturates and weakens at higher intensities of the guided pulse.

Since the focusing strength is determined by the curvature of the χ_0 rather than by its absolute value, a possible solution to the nonlinear saturation of the EIG is to use a first-order Bessel beam, which has an intensity node on axis. The density perturbation on axis vanishes while its curvature does not. To utilize the J_1 Bessel beam, frequency detuning has to be chosen $\Delta\omega > \omega_p, \tau_L^{-1}$. For a guided beam of arbitrary longitudinal shape $\chi_0 \approx 2|a_1|^2 \omega_0^2/(\omega_p^2 - \Delta\omega^2)$, enabling focusing by a beam with intensity minimum on axis. The guiding condition similar to Eq. (8) can be derived for the $J_1(r/\sigma_1)$ beam. The strongest focusing occurs for $\sigma_1 = 1.04\sigma$, yielding $k_0^2\sigma^2\alpha_1^2 = 5[(\Delta\omega/\omega_p)^2 - 1]$, from which the power threshold condition can be derived:

$$P_{B1} = 3.6 \ \ \mathrm{GW}(\Delta\omega^2/\omega_p^2 - 1)\frac{L_g}{k_0\sigma^2}. \tag{10}$$

The linear theory breaks down at $r \approx \sigma$ if $\alpha_0 > 2.7\lambda_0\sigma/\lambda_p^2$. In practice, this limitation is likely to be overstringent since the nonlinear saturation of the EIG in a limited spatial region may not strongly influence the overall focusing. As a numerical example, consider guiding a $\tau_L = 10$ fs long, $\sigma = 18\,\mu$m wide laser pulse through 1 cm of $n_0 = 10^{19}$ cm^{-3} plasma. For $\Delta\omega/\omega_p = 1.5$ the threshold power of the Bessel beam $P_{B1} \approx 16$ GW. This beam can guide a pulse with a normalized vector potential up to at least $\alpha_0 = 0.5$, or 2.2 TW.

The guided pulse generates an appreciable long-wavelength ($\lambda = \lambda_p$) plasma wake, which can interfere with the short-wavelength plasma wake responsible for guiding [27]. If the fractional density modulation of the fast (long-wavelength) wake is $\hat{n}_2$, then, according to Ref. [27], the fast wake starts suppressing the slow wake when $n^2 k_0/k_p \approx 1$. This corresponds roughly, to the parameters above. The physical mechanism behind such suppression is that an electron, executing plasma oscillations in the field of the fast wake, sweeps across a wavelength of the slow wake. However, this effect is likely to be small for very short (of order ω_p^{-1}) pulses, since there is not enough time for the electrons to get sufficiently displaced. However, guiding suppression due to the fast wake generation deserves careful consideration.

V. CONCLUSIONS

In conclusion, we demonstrated analytically and numerically the electromagnetically induced guiding of two counterpropagating lasers in the plasma at intensities much below the threshold for relativistic guiding. In addition, we described the technique for electromagnetic channeling of an ultrashort laser pulse by a lower-intensity counterpropagating Bessel beam. Depending on the frequency detuning between the guided and guiding beams, the zeroth- or first-order Bessel beams can be used for guiding ultrashort pulses. A first-order Bessel beam requires higher power, but enables guiding higher intensity pulses.

In this work we assumed that the two lasers are detuned sufficiently far from the plasma resonance, so that the interaction between the beams is purely refractive. Investigating possible energy exchange between the beams is the subject of the future work.

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