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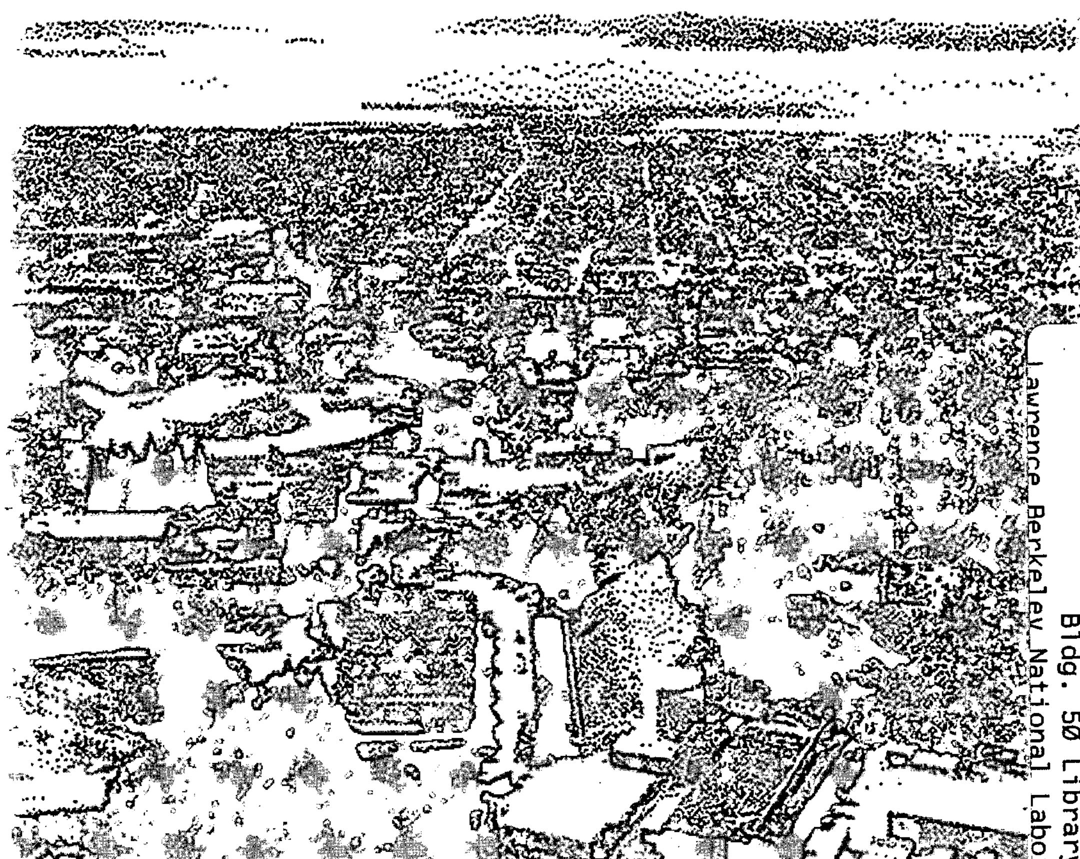

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Laser Acceleration in Plasmas with Capillary Waveguide

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(February 11, 1999)

We present new methods for ponderomotive-driven laser acceleration in plasmas over much extended distances, many meters long in a single stage. The diffraction limit on acceleration distance is overcome by enclosing a column of uniform plasma with a capillary waveguide. For laser wakefield acceleration in particular, we provide a full 3D solution of wakefield in the waveguide, and propose two methods to overcome the phase slippage limit as well. Effect of nonuniform plasma on waveguide performance is also analyzed.

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I. INTRODUCTION

Given the extensive exposition of capillary waveguide for direct-field acceleration in vacuum and gases [1], it would appear natural to apply the same waveguide for ponderomotive-driven acceleration in plasmas. Nevertheless, doing so represents a radical conceptual deviation from the current mainstream development of plasma-based laser acceleration [2]. First of all, the prevailing notion on optical guiding in plasmas is based on an analogy to optical fiber in which the index of refraction is maximum on the axis, opposite to that of capillary waveguide. Secondly, it is commonly believed that no solid-state waveguide structure could sustain the laser power required for acceleration in plasmas without being, at least partially, turned into plasmas. Therefore, despite the proven capability of capillary waveguide in guiding lasers of both high average and peak powers in vacuum and gases, the only method under active investigation for optical guiding in plasmas is to tailor the plasma itself in transverse density profile one way or the other, by either relativistic self-focusing [3], hydrodynamic expansion [4], capillary discharge [5], or hollow plasma channel [6].

However, all methods relying on dynamic focusing from plasmas rather than passive guiding from durable external structures suffer from severe limitations of the approach. In the case of relativistic self-focusing, a laser pulse has to make its own channel for guiding. Therefore, it is subject to various laser-plasma instabilities in the long pulse regime and ineffective for guiding short pulses [7]. With hydrodynamic expansion, on the other hand, the waveguide is generated by another laser which itself needs to be guided. In absence of such a pre-required waveguide for the driver laser, a line focus from an axicon is used instead. Thus the length of the waveguide so produced is limited by what can be provided with an axicon in the first place. In a capillary discharge, a nonuniform plasma is created by ablating wall materials with high discharge current. As such, this approach is impractical for accelerator applications because of its short usable lifetime. Lastly, the hollow plasma channel, envisioned as a vacuum core embedded in a uniform plasma, is only a fictitious idea. As a result, optical guiding in plasmas has been demonstrated over distances only up to a few centimeters. Although impressive when measured against the extremely short Rayleigh length under the experimental condition, such a waveguide performance in terms of real distance is still far from being practical.

In this letter, we propose to solve the problem of optical guiding in plasmas by a durable solid-state waveguide. In particular we consider laser wakefield acceleration [8] taken place in a capillary waveguide with a core of uniform plasma. Furthermore, two methods are provided to sustain acceleration over multiple phase slippage lengths, thus many meters long in a single stage. Effect of nonuniform plasma core on waveguide performance is also analyzed. The notations used here follow closely that introduced in the companion letter [1].

II. LASER WAKEFIELD ACCELERATION

Wave equation for laser field propagation in a weakly relativistic plasma is governed by [9]

$$\left[\nabla^2 - \frac{1}{c^2}\frac{\partial^2}{\partial t^2}\right]\mathbf{E}_l = \frac{\omega_p^2}{c^2}\left[1+\frac{\delta n}{n_0}-\frac{a^2}{2}\right]\mathbf{E}_l \ , \qquad (1)$$

where the plasma density modulation $\delta n/n_0$, while driven by the ponderomotive potential of a laser pulse $a^2{=}<|e\mathbf{E}_l/mc\omega|^2>$, could generate a wakefield $\mathbf{E}_w = -\nabla\Phi$, through Poisson's equation $\nabla^2\Phi = (e/\epsilon_0)\delta n$, where the angle bracket indicates time average over an optical period. For a cold plasma fluid the wake potential is determined in the weakly relativistic limit by [10]

$$\left[\frac{\partial^2}{\partial t^2}+\omega_p^2\right]\Phi = \omega_p^2\frac{mc^2}{e}\frac{a^2}{2} \ . \qquad (2)$$

Under the condition $a^2 \ll 1$, we have $\delta n/n_0 \ll 1$, as will be shown later. As a result, the second and third term on the right of Eq.(1) can be dropped and the wave equation is then decoupled from the plasma equation. In this case, the only effect of a plasma on laser propagation is through its index of refraction.

The phase velocity of a wakefield is same as the group velocity of the driver laser pulse. For EH_{11} mode, the group velocity is $v_g = d\omega/d\beta_{11} = c(1 - 1/2\gamma_g^2 - 1/2\gamma_p^2)$, corresponding to an acceleration phase slippage length

$$L_a = \frac{\lambda_p}{|1/\gamma_g^2 + 1/\gamma_p^2 - 1/\gamma^2|} \ . \tag{3}$$

To derive laser wakefield and its characteristic properties, we take an approach similar to a 1D linear analysis [10], except here the solution we provide is in full 3D. Introducing a variable $\zeta = z - v_g t$, Eq.(2) is solved as

$$\Phi = -(k_p mc^2/e) \int_\zeta^\infty d\zeta' \sin[k_p(\zeta - \zeta')] \frac{a^2}{2} \ , \tag{4}$$

where $k_p = \omega_p/v_g$. For a Gaussian pulse of EH_{11} mode

$$a^2(\rho,\zeta) = \frac{a_0^2}{2} J_0^2(U_{11}\rho) e^{-\zeta^2/2\sigma_z^2 - z/L_{attn}} \ , \tag{5}$$

where $\rho = r/R$, wake potential behind a laser pulse is

$$\begin{aligned} \Phi &= -\Phi_0 J_0^2(U_{11}\rho) e^{-z/L_{attn}} \sin(k_p z - \omega_p t) \ , \\ \Phi_0 &= (\sqrt{2\pi} mc^2/4e) a_0^2 k_p \sigma_z e^{-(k_p\sigma_z)^2/2} \ . \end{aligned} \tag{6}$$

The longitudinal wakefield is then given by

$$E_{wz} = E_a J_0^2(U_{11}\rho) e^{-z/L_{attn}} \cos(k_p z - \omega_p t) \ , \tag{7}$$

and the transverse wakefield by

$$\begin{aligned} E_{wr} = -2(\gamma_p/\gamma_g) E_a J_0(U_{11}\rho) J_1(U_{11}\rho) \\ e^{-z/L_{attn}} \sin(k_p z - \omega_p t) \ , \end{aligned} \tag{8}$$

where the peak acceleration field, $E_a = \Phi_0 k_p$, is maximized if the laser pulse length is chosen according to the condition, $k_p\sigma_z = 1$. From here on, this condition will be used unless otherwise stated. Indeed, with

$$\begin{aligned} \frac{\delta n}{n_0} = \frac{\sqrt{\pi} a_0^2}{\sqrt{8\exp(1)}} \left\{ 1 + \frac{2\gamma_p^2}{\gamma_g^2} \left[1 - \frac{J_1^2(U_{11}\rho)}{J_0^2(U_{11}\rho)} \right] \right\} \\ J_0^2(U_{11}\rho) e^{-z/L_{attn}} \sin(k_p z - \omega_p t) \ . \end{aligned} \tag{9}$$

we have $\delta n/n_0 \ll 1$, if $a_0^2 \ll 1$.

As wakefield is excited in a plasma channel, energy in the driver pulse is depleted. A characteristic pump depletion length can be defined by the condition, $W_l = W_w$, where W_l is the initial energy of the laser pulse

$$W_l = \pi\sqrt{\pi/8} J_1^2(U_{11}) mc^2 a_0^2 \gamma_p^2 R^2 / r_e \lambda_p \ , \tag{10}$$

and W_w is the energy in the wakefield the laser pulse left behind as it propagates a distance of L_{pump}

$$W_w = \frac{\pi^3 mc^2 a_0^4 R^2 L_{pump}}{16\exp(1)\lambda_p^2 r_e} \left[I_z + (\gamma_p/\gamma_g)^2 I_r \right] \ . \tag{11}$$

The two terms above on the right correspond to energy in the longitudinal and transverse wakefield, respectively, where $I_z = \int_0^1 d\rho \rho J_0^4(U_{11}\rho) = 0.0762$ and $I_r = 4\int_0^1 d\rho\rho J_0^2(U_{11}\rho) J_1^2(U_{11}\rho) = 0.102$. Hence, we have

$$L_{pump} = \frac{4\sqrt{2\pi}\exp(1) J_1^2(U_{11}) \gamma_p^2 \lambda_p}{\pi^2 a_0^2 [I_z + (\gamma_p/\gamma_g)^2 I_r]} \ . \tag{12}$$

In addition, we define a characteristic pulse dispersion length over which the driver pulse double its length

$$L_{disp} = \frac{\sqrt{3}\gamma_p^4 \lambda}{\pi[1 + (\gamma_p/\gamma_g)^2]} \ . \tag{13}$$

Finally, the beta function due to transverse wakefield is

$$\beta_t = (2/U_{11})[\exp(1)/2\pi]^{1/4} \sqrt{\frac{\gamma}{\sin\phi_a} \frac{R}{a_0}} \ . \tag{14}$$

III. SOLUTIONS TO PHASE SLIPPAGE

Two methods are proposed in this section to overcome the limit on acceleration distance set by the phase slippage length of Eq.(3). The first method requires inserting plasma layers of higher density, each of length L_d, as drift sections in between acceleration sections, each of length L_a. Two conditions need to be satisfied for this method. First, the length of a drift section is determined by

$$L_d = \frac{\lambda_p}{1/\gamma_g^2 + 1/\gamma_{pd}^2 - 1/\gamma^2} \ , \tag{15}$$

where $\gamma_{pd} = \lambda_{pd}/\lambda$ and λ_{pd} is the plasma period corresponding to the plasma density in the drift section. This condition guarantees energy gain over the full length in each acceleration section, since L_d is the distance for the particle to slip π phase with respect to the acceleration wave of period λ_p. Secondly, the plasma density in the drift section is set according to $\lambda_p/\lambda_{pd} = 2m$, where m is an integer. This condition ensures that there is no net energy exchange between the particle and the laser wakefield excited in the drift section, since L_d is also the distance over which a particle slips $2m\pi$ phase with respect to the wakefield with period λ_{pd} in the drift section. Thus the average gradient that can be maintained over multiple slippage lengths is $G = \Delta W_a/(L_a + L_d)$, where $\Delta W_a = eE_a L_a T_a$ is the energy gain in one acceleration section. In a limiting case with $(\gamma/\gamma_p)^2 \gg 1$ and $(\gamma_g/\gamma_p)^2 \gg 1$, we have $L_a/L_d = (\lambda_p/\lambda_{pd})^2$. It is noted that the ratio λ_p/λ_{pd} does not have to be exactly an even integer, as it can be shown through Eq.(7) that any residual energy exchange occurring in a drift section would be small if the driver pulse length is matched to the plasma period in the acceleration section. Only in one special case when $\lambda_p/\lambda_{pd} = 1$, energy gain in an acceleration section is completely canceled by energy loss due to deceleration in a drift section. An example with drift section approach is given in Table 1 for highly relativistic electron satisfying $(\gamma/\gamma_p)^2 \gg 1$ and $(\gamma/\gamma_g)^2 \gg 1$.

Table 1. Example with Drift Section.

λ [μm]	1	n_0 [10^{17}/cm^3]	1.1	E_a [GV/m]	0.94
R/λ	150	$(\delta n/n_0)_{\max}$	0.033	E_s [GV/m]	1.7
ν_2	1.5	γ_p	100	L_a [m]	0.94
P_0 [TW]	20	γ_g	392	L_d [m]	0.25
W_l [J]	2.7	σ_z [μm]	16	L_{attn} [m]	7.9
a_0	0.28	ΔW_a [MeV]	560	L_{disp} [m]	52
λ_p/λ_{pd}	2	G [GeV/m]	0.47	L_{pump} [m]	117

The second method utilizes longitudinal modulation in laser intensity due to beating of two waveguide modes. The idea is to choose the beating period same as the distance for a 2π phase slippage, such that the wakefield is stronger when the particle is in accelerating phase, and weaker in decelerating phase, resulting in net energy gain over multiple slippage lengths. When two modes are included, Eq.(5) is modified to

$$a^2(\rho,\zeta,z) = \frac{a_0^2}{2}|f_b(\rho,z)|^2 e^{-\zeta^2/2\sigma_z^2} \ , \tag{16}$$

where the profile, normalized to $f_b(0,0)=1$, is given by

$$\begin{aligned} f_b(\rho,z) &= \tfrac{1}{1+\eta}\left[E_{11}(\rho,z)+\eta E_{12}(\rho,z)\right] \ , \\ E_{11}(\rho,z) &= J_0(U_{11}\rho)e^{i\beta_{11}z-\alpha_{11}z} \ , \\ E_{12}(\rho,z) &= J_0(U_{12}\rho)e^{i\beta_{12}z-\alpha_{12}z} \ . \end{aligned} \tag{17}$$

Assuming $(\gamma_{g11}/\gamma_p)^2 \gg 1$, $(\gamma_{g12}/\gamma_p)^2 \gg 1$ and $(\gamma/\gamma_p)^2 \gg 1$, where γ_{g11} and γ_{g12} are γ_g factors for EH_{11} and EH_{12} modes, respectively, the group velocity and slippage length then become same for both modes, i.e., $v_g = c(1-1/2\gamma_p^2)$ and $L_a = \gamma_p^2\lambda_p$. There are three characteristic length scales: $l_1 \sim \{1/\alpha_{11}, 1/\alpha_{12}\}$ is due to mode attenuation; $l_2 \sim 1/(\beta_{11}-\beta_{12})$ is due to beating of the two modes; and $l_3 \sim 1/k_p$ is the plasma period. As they satisfy $l_1 \gg l_2 \gg l_3$, we deduce from Eq.(4) that

$$E_{wz} = E_a|f_b(\rho,z)|^2\cos(k_pz-\omega_pt) \ , \tag{18}$$

and, in particular, for acceleration field on the axis

$$|f_b(0,z)|^2 = \frac{1+\eta^2+2\eta\cos[(\beta_{11}-\beta_{12})z]}{(1+\eta)^2} \ . \tag{19}$$

By requiring $\beta_{11}-\beta_{12} = \pi/L_a$, we have the condition

$$\gamma_{g11}^2/\gamma_p^3 = U_{12}^2/U_{11}^2 - 1 \ , \tag{20}$$

where $U_{12} = 5.52$. Energy gain over $2L_a$ distance is then $\Delta W_{2\pi} = \int_0^{2L_a} eE_{wz}dz = eE_a2L_aT_{2\pi}$, where $T_{2\pi} = \eta/(1+\eta)^2$. As expected, $T_{2\pi}$ vanishes when there is only one mode with $\eta=0$, and it reaches a maximum when the two modes have equal amplitude on the axis with $\eta = 1$. The relative mode amplitude can be adjusted easily by changing TEM_{00} mode waist w_0 at waveguide entrance [11] according to this relation

$$\eta = \frac{J_1^2(U_{11})\int_0^1 J_0(U_{12}\rho)\exp[-\frac{\rho^2}{(w_0/R)^2}]\rho d\rho}{J_1^2(U_{12})\int_0^1 J_0(U_{11}\rho)\exp[-\frac{\rho^2}{(w_0/R)^2}]\rho d\rho} \ . \tag{21}$$

An example with mode beating is given in Table 2 for highly relativistic electron. Here the average gradient is defined by $G = \Delta W_{2\pi}/2L_a$. It is interesting to note that 99.86% of power from the free-space TEM_{00} mode is coupled into the two dominant waveguide modes, 90.17% in EH_{11} mode, 9.69% in EH_{12} mode, 0.11% in all other modes, and only 0.03% misses the core of the waveguide.

Table 2. Example with Mode Beating.

λ [μm]	1	n_0 [10^{17}/cm^3]	4.1	$\Delta W_{2\pi}$ [MeV]	57
R/λ	300	γ_p	52	G [GeV/m]	0.2
ν_2	1.5	w_0/R	0.496	L_a [m]	0.14
σ_z [μm]	8.3	η	0.5	P_0 [TW]	40

IV. EFFECT OF NONUNIFORM PLASMA

We have analyzed the capillary waveguide when the core is filled with a plasma of uniform density n_0. In reality this can only be an approximation. To evaluate the effect of a nonuniform plasma, let us consider a special case when the uniform background is modified by a parabolic profile, $n = n_0 + \Delta n[1-(r/R)^2]$. Such a profile has a defocusing effect on the mode if $\Delta n > 0$, thus making the waveguide less effective. However, the attenuation length for the EH_{11} mode is reduced by a factor of two at most, if the following criteria is satisfied

$$\frac{\Delta n}{n_0} \le 4.6\,\frac{\gamma_p^2}{\gamma_g^2} \ . \tag{22}$$

For the example in Table 1, this corresponds to $\Delta n/n_0 \le$ 30%. Thus we may infer from here that the guiding provided by the capillary waveguide can also be rather stable, against either systematic or random variations in plasma density. In the remainder of this section, we will first establish a general-purpose perturbation technique for the study of capillary waveguide with a transversely nonuniform core profile, and then derive the criteria, Eq.(22), as a special case for EH_{11} mode with the parabolic core profile.

It is noted that a similar perturbation technique can be derived from the scalar wave equation for bound modes under weak-guidance approximation, $|(\nu_2-\nu_1)/\nu_1| \ll 1$ [12]. However, this result can not be used here without justification, as the approximation does not apply. Instead, we have to start from the vector wave equation

$$[\nabla_t^2 + \nu^2 k^2 - \Gamma^2]\mathbf{E}_t = -\nabla_t(\mathbf{E}_t \cdot \nabla_t \ln \nu^2) ,$$
$$[\nabla_t^2 + \bar{\nu}^2 k^2 - \bar{\Gamma}^2]\bar{\mathbf{E}}_t = -\nabla_t(\bar{\mathbf{E}}_t \cdot \nabla_t \ln \bar{\nu}^2) , \quad (23)$$

for both perturbed and unperturbed waveguides, to take into account the abrupt change in index of refraction at the core-cladding interface. From Eq.(23), we obtain

$$(\Gamma^2 - \bar{\Gamma}^2) < \mathbf{E}_t \cdot \bar{\mathbf{E}}_t > = k^2 < (\nu^2 - \bar{\nu}^2)\mathbf{E}_t \cdot \bar{\mathbf{E}}_t > + < \bar{\mathbf{E}}_t \cdot \nabla_t^2 \mathbf{E}_t - \mathbf{E}_t \cdot \nabla_t^2 \bar{\mathbf{E}}_t > + < \bar{\mathbf{E}}_t \cdot \nabla_t(\mathbf{E}_t \cdot \nabla_t \ln \nu^2) - \mathbf{E}_t \cdot \nabla_t(\bar{\mathbf{E}}_t \cdot \nabla_t \ln \bar{\nu}^2) > , \quad (24)$$

where the angle bracket denotes integration over cross section, $\int_{A_\infty} dA = \int_{core} dA + \int_{cladding} dA$. As discussed before [1] for oversized waveguide, all modes have the radial dependence $e^{ik_1 r\sqrt{\bar{\nu}^2-1}}/\sqrt{r}$ for $r \geq R$, which would make the integral over the cladding infinite if $\hat{\nu}$ is real, as in the case for a radiation mode. This difficulty of purely mathematical origin can be avoided by introducing a positive imaginary part of proper magnitude to ν_2, corresponding to an artificial dielectric loss in the cladding. It can be shown that such a procedure would not affect either mode profile in the core or the propagation constant, provided that $\mathrm{Im}(\nu_2) \ll \mathrm{Re}(\nu_2)$. As a result, the integral over the cladding is negligible and the second term on the right of Eq.(24) vanishes. To treat the third term, we further specify the profiles as

$$\bar{\nu}^2 = \begin{cases} \nu_1^2 + \bar{\mu}(r) & : \ r \leq R \\ \nu_2^2 & : \ r > R , \end{cases}$$
$$\nu^2 = \begin{cases} \nu_1^2 + \mu(r) & : \ r \leq R \\ \nu_2^2 & : \ r > R , \end{cases} \quad (25)$$

and, without loss of generality, require that $\mu(r=R) = \bar{\mu}(r=R) = 0$. Therefore, $\ln \nu^2 = \ln(\nu_1^2 + \mu) + [\ln \nu_2^2 - \ln(\nu_1^2 + \mu)]H(r-R)$, where $H(r-R)$ is a step function, and, $\nabla_t \ln \nu^2 = \nabla_t \ln \bar{\nu}^2 \cong \ln(\nu_2^2/\nu_1^2)\ \delta(r-R)\hat{\mathbf{r}}$, where $\hat{\mathbf{r}}$ is a unit radial vector. Assuming $|\nu^2 - \bar{\nu}^2|/\bar{\nu}^2 \ll 1$, a first order perturbation for Γ can be obtained by substituting zeroth order solution for the eigenfunction $\mathbf{E}_t \cong \bar{\mathbf{E}}_t$. The third term thus also vanishes. Finally, we have

$$\Gamma^2 - \bar{\Gamma}^2 = \frac{k^2 \int_0^R \int_0^{2\pi} (\nu^2 - \bar{\nu}^2) \bar{\mathbf{E}}_t^2 r dr d\phi}{\int_0^R \int_0^{2\pi} \bar{\mathbf{E}}_t^2 r dr d\phi} . \quad (26)$$

We now apply Eq.(26) to EH_{11} mode with $\bar{\mu}(r) = 0$ for unperturbed waveguide and $\mu(r) = \Omega[(r/R)^2 - 1]$ for perturbed waveguide. Using the mode profile $\bar{E}_y = E_0 J_0(k_{r1} r)$, the definition $k_{r1} = (U_{11} - i\Lambda/\gamma_g)/R$, and the expansion $J_0(k_{r1} r) = J_0(U_{11}\rho) + i(\Lambda/\gamma_g)\rho J_1(U_{11}\rho)$, we obtain to leading order

$$\Gamma = \bar{\Gamma} - I_A k\Omega/2 + iI_B k\Omega\Lambda/\gamma_g , \quad (27)$$

where $I_A = 1 - I_3/I_1, I_B = I_4/I_1 - I_2 I_3/I_1^2, I_1 = \int_0^1 J_0^2(U_{11}\rho)\rho d\rho = 0.135, I_2 = \int_0^1 J_0(U_{11}\rho)J_1(U_{11}\rho)\rho^2 d\rho = 0.056, I_3 = \int_0^1 J_0^2(U_{11}\rho)\rho^3 d\rho = 0.0294$, and $I_4 = \int_0^1 J_0(U_{11}\rho)J_1(U_{11}\rho)\rho^4 d\rho = 0.0244$. By requiring the attenuation rate for the unperturbed mode is not more than doubled due to the waveguide perturbation, we have $\Omega \leq 1/I_B U_{11}\gamma_g^2$. This then leads to the criteria, Eq.(22), upon using the relation, $\Omega = \Delta n/n_0\gamma_p^2$.

V. CONCLUSIONS

I have introduced a new paradigm for plasma-based laser acceleration by replacing dynamic focusing from plasmas with passive guiding from durable solid-state waveguides. With the new approach, acceleration distance can be increased dramatically. As a result, various detrimental laser-plasma instabilities can be avoided as well by choosing appropriate regime of operation, in particular, with $a_0^2 \ll 1$, $\delta n/n \ll 1$, $P_0/P_c < 1$ and $k_p\sigma_z = 1$, where $P_c[\mathrm{GW}] = 17\gamma_p^2$ is the critical power for relativistic focusing [7]. In contrast, the current development of plasma-based laser acceleration is trapped in a dilemma due to the lack of effective means for optical guiding. To achieve respectable energy gain over severely limited distance, acceleration gradient has to be driven up either close to wavebreaking limit or well into the regime prone to various laser-plasma instabilities [13]. In return, further limitations are imposed on acceleration distance, electron beam quality and system controllability.

Finally, it is noted that most principles and techniques introduced here can be applied to other types of ponderomotive-driven laser acceleration. In particular, extension to plasma beatwave acceleration [8] is straightforward: As for the creation of plasma column, ionization of gases by a laser propagating in the waveguide appears to be both effective and efficient. Furthermore, the open iris-loaded waveguide [14] can also be used with the acceleration schemes presented here. These and other related subjects will be treated in a forthcoming series of papers. This work was supported by the U.S. Department of Energy under contract No.DE-AC03-76SF00098.

[1] M. Xie, *Laser Acceleration in Vacuum and Gases with Capillary Waveguide*, LBNL-42783, submitted to *Phys. Rev. Lett.* on February 11, 1999.
[2] *AIP Conf. Proc.*, **335** (1995), **398** (1997), **472** (1999).
[3] A. Borisov et al., *Phys. Rev. Lett.*, **68**, 2309(1992).
[4] C. Durfee et al., *Phys. Rev. Lett.*, **71**, 2409(1993).
[5] Y. Ehrlich et al., *Phys. Rev. Lett.*, **77**, 4186(1996).
[6] T. Chiou et al., *Phys. Plasmas*, **2**, 310(1995).
[7] E. Esarey et al., *IEEE Tran. Plasma Sci.*, **24**, 252(1996).
[8] T. Tajima et al., *Phys. Rev. Lett.*, **43**, 267(1979).
[9] T. Antonsen et al., *Phys. Fluids B*, **5**, 1440(1993).
[10] L. Gorbunov et al., *Sov. Phys. JETP*, **66**, 290(1987).
[11] R. Jenkins et al., *Applied Optics*, **31**, 5086(1992).
[12] A. Snyder et al., *J. Opt. Soc. Am.*, **72**, 1131(1982).
[13] T. Katsouleas, *AIP Conf. Proc.*, **398**, 175(1997).
[14] M. Xie, *Proc. of 1997 Part. Accel. Conf.*, 660 (1997).

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