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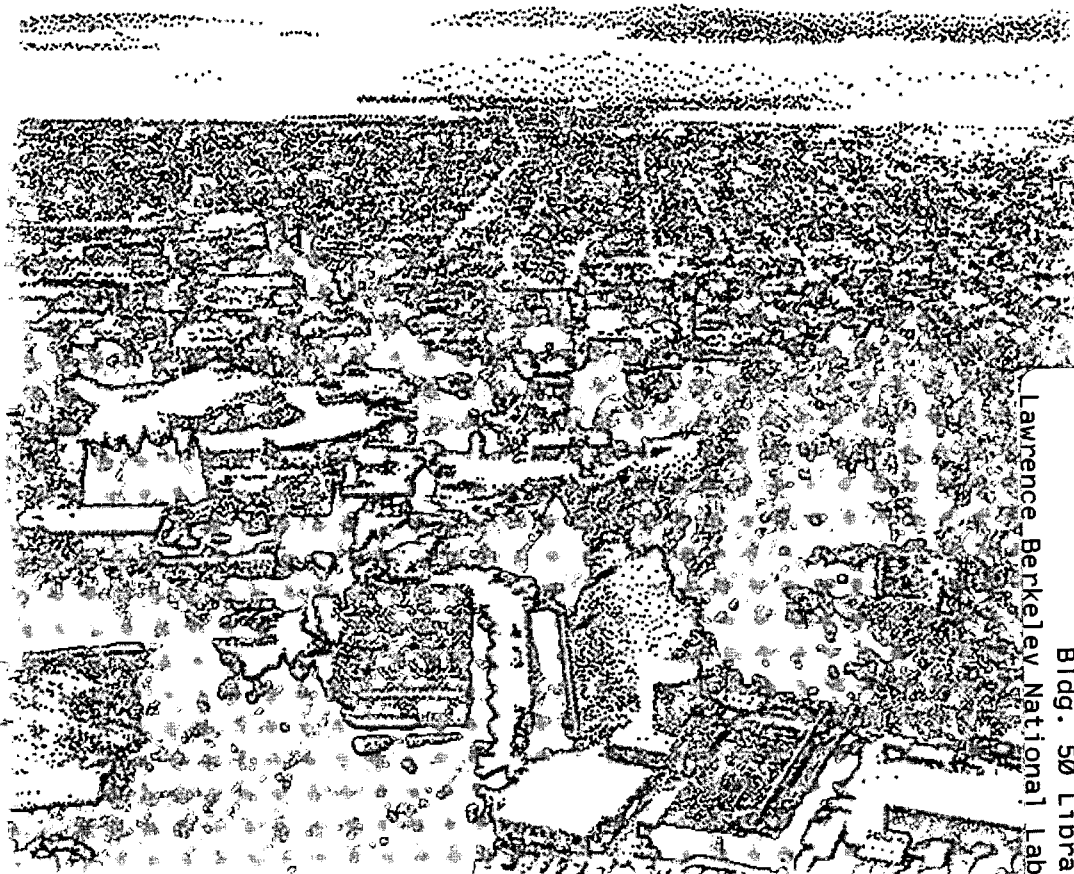
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Accelerator and Fusion
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(February 11, 1999)

A unified framework is developed to overcome simultaneously three major limitations on acceleration distance and hence on the feasibility of two classes of laser acceleration. The three limitations are due to laser diffraction, acceleration phase slippage, and structure damage by high power laser if solid-state optical waveguide is used. The two classes of laser acceleration are direct-field acceleration and ponderomotive-driven acceleration. This letter and its companion [1] provide solutions that are crucial to all mainstream approaches for laser acceleration, either in vacuum, gases or plasmas.

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I. INTRODUCTION

The tremendous growth in accelerator energy of this century has been truly remarkable. However, the workhorse behind this success, microwave powered structure, has more or less reached its known limit in acceleration gradient. To continue the role well into next century as a primary instrument for exploration of fundamental physics in high energy frontier, accelerator has to be powered by electromagnetic (EM) wave of much shorter wavelength. Scaling current technology such as structure dimension down with wavelength is not an attractive pass, at least in the long run. The demand for high gradient acceleration has stimulated widespread interest and research on laser accelerations that operate under fundamentally different concepts and principles.

Yet, despite the impressive progress made in recent years [2], high gradient has been demonstrated only over very short acceleration distances, up to a couple of centimeters. Standing in between proof of principles and real-world applications are three major limitations [3–10]. First of all, the distance of laser-particle interaction is limited by laser diffraction. Secondly, the distance over which a particle can continuously gain energy from a wave is limited by acceleration phase slippage due to the difference between particle speed and phase velocity of the wave. Finally, an acceleration structure, if capable of both guiding laser field and harnessing energy transfer, has to withstand high power laser, and its effectiveness is limited by laser damage threshold. Evidently, the key entry for a breakthrough is at the last point.

In this letter and its companion [1] we present a unified framework to overcome all three major limitations on laser acceleration. Based on interaction mechanism,

laser acceleration can be divided into two classes: direct acceleration with the longitudinal field, and acceleration driven by ponderomotive force of the transverse field of a laser. Accordingly, the two classes will be treated separately by the two letters. However, the following discussion on optical waveguide applies to both.

II. PROPERTIES OF WAVEGUIDE MODES

The capillary waveguide considered here is made of a hollow core with an index of refraction ν_1 and radius R , embedded in a medium of dielectric or metal with a complex index of refraction ν_2 . We are interested in oversized waveguide satisfying $\lambda_1/R \ll 1$, where $\lambda_1 = \lambda/\nu_1$ and λ is the wavelength in vacuum. As a result, EM wave in the core is dominantly transverse. Capillary waveguide of this type has been used widely for industrial and medical applications to transport lasers of high average power [11]. More recently, it has also been tested at high peak power [12,13]. In this section we explore and elaborate the mode properties important for laser acceleration.

The eigenmodes of the waveguide can be solved following the same procedure by Marcatili et al. [14] under the condition $|\sqrt{\hat{\nu}^2 - 1}| \gg \lambda_1/R$, where $\hat{\nu} = \nu_2/\nu_1$. Expressing the eigenmodes in the following form

$$\begin{Bmatrix} \mathcal{E}(r, \phi, z, t) \\ \mathcal{H}(r, \phi, z, t) \end{Bmatrix} = \begin{Bmatrix} \mathbf{E}_{lm}(r, \phi) \\ \mathbf{H}_{lm}(r, \phi) \end{Bmatrix} e^{i(\beta_{lm}z - \omega t) - \alpha_{lm}z}, \quad (1)$$

the eigenvalues are given by

$$\beta_{lm} = k_1(1 - 1/2\gamma_g^2), \quad \alpha_{lm} = \text{Re}(\Lambda)/\gamma_g^2 R, \quad (2)$$

where $k_1 = \nu_1 k$, $k = 2\pi/\lambda$, $\gamma_g = 2\pi R/U_{lm}\lambda_1 \gg 1$, and U_{lm} is the m th root of the equation $J_{l-1}(U_{lm}) = 0$. There are three types of modes, corresponding to

$$\Lambda = \begin{cases} \frac{1}{\sqrt{\hat{\nu}^2 - 1}} & : TE_{0m} \ (l=0) \\ \frac{\hat{\nu}^2}{\sqrt{\hat{\nu}^2 - 1}} & : TM_{0m} \ (l=0) \\ \frac{\hat{\nu}^2 + 1}{2\sqrt{\hat{\nu}^2 - 1}} & : EH_{lm} \ (l \neq 0) \end{cases} \quad (3)$$

For laser acceleration, we are interested primarily in two low-order modes: TM_{01} mode for direct-field acceleration in vacuum and gases with its on-axis E_z component, and EH_{11} mode for ponderomotive-driven acceleration in plasmas. Accordingly, we consider three cases: $\delta\nu_1 = 0$ when the core is in vacuum, $\delta\nu_1 > 0$ and $\delta\nu_1 < 0$ when the core is filled with uniform gases and plasmas, respectively, where $\delta\nu_1 = \nu_1 - 1$ and $|\delta\nu_1| \ll 1$.

The electric fields within the core $r \leq R$ are given by

$$TM_{01} : \begin{cases} E_z = E_a J_0(k_{r1}r) \\ E_r = -i(\Gamma/k_{r1})E_a J_1(k_{r1}r), \end{cases} \quad (4)$$

$$EH_{11} : \begin{cases} E_y = E_0 J_0(k_{r1}r) \\ E_z = -i(k_{r1}/\Gamma)E_0 J_1(k_{r1}r) \sin \phi, \end{cases} \quad (5)$$

where E_a is the peak acceleration field for TM_{01} mode, E_0 is the peak transverse field for EH_{11} mode, $\Gamma = \beta_{lm} + i\alpha_{lm}$ and $k_{r1} = (U_{lm} - i\Lambda/\gamma_g)/R$. To leading order $\Gamma/k_{r1} = \gamma_g$, since by definition $k_{r1}^2 = \Gamma^2 + k_{r1}^2$. Given electric field, magnetic field of a mode can be determined by $\mathbf{H}_t = \hat{\mathbf{z}} \times (\Gamma \mathbf{E}_t + i\nabla_t E_z)/kZ_0$, $H_z = (i/\Gamma)\nabla_t \cdot \mathbf{H}_t$, where subscript t denotes transverse component of a vector or operator, $\hat{\mathbf{z}}$ is a unit vector in z -direction, and Z_0 is the vacuum impedance. As seen from Eqs.(4,5), the transverse fields dominate over the longitudinal ones by a large factor γ_g . For TM_{01} mode, E_r is peaked at $r/R = 0.481$ with a maximum value $E_r^{max} = 0.582\gamma_g E_a$. For EH_{11} mode, E_y is peaked on the axis. To evaluate surface field E_s at $r = R$, we expand Bessel functions in the transverse fields of Eqs.(4,5) using the expression for k_{r1} and keeping the larger one of the two components

$$\begin{aligned} E_s/E_a &= \max\{1, |\Lambda|\} |J_0(U_{01})| & : TM_{01}, \\ E_s/E_0 &= \max\{1, |\Lambda|\} |J_1(U_{11})|/\gamma_g & : EH_{11}. \end{aligned} \quad (6)$$

Here we come upon one of the most important advantages of the capillary waveguide: for TM_{01} mode, surface field can be smaller than peak acceleration field, superior to other schemes [4–10] including even microwave structure; and for both modes, surface fields are much smaller than peak transverse fields. According to experimental data on dielectrics at $\lambda = 1\mu\text{m}$ [15], damage threshold in terms of surface field is around $5 \sim 10$ GV/m for pulse length of 1 ps or shorter, thus making acceleration in GV/m scale possible. Within the cladding $r \geq R$, all field components have the radial dependence $e^{ik_{r2}r}/\sqrt{r}$, where to leading order $k_{r2} = k_1\sqrt{\hat{\nu}^2 - 1}$. A non-vanishing imaginary part of $\hat{\nu}$ due to dielectric loss would give rise to exponential decay in radial direction. Hence power in each mode is distributed dominantly within the core and it can be expressed as $P(z) = P_0 e^{-z/L_{attn}}$, where $L_{attn} = 1/2\alpha_{lm}$ is power attenuation length and to leading order

$$P_0 = \begin{cases} \pi R^2 \gamma_g^2 E_a^2 J_0(U_{01})^2 / 2Z_0 & : TM_{01} \\ \pi R^2 E_0^2 J_1(U_{11})^2 / 2Z_0 & : EH_{11}. \end{cases} \quad (7)$$

It is noted that EH_{11} mode is linearly polarized, whereas TM_{01} mode is radially polarized. A linearly polarized mode with non-vanishing on-axis E_z can be constructed by a proper mixing of TM_{01} with EH_{21} mode [16]. Electric field for the mixed mode is given by

$$\begin{aligned} E_z &= E_a [J_0(k_{r1}r) + J_2(k_{r1}r) \cos 2\phi] \\ E_y &= -2i(\Gamma/k_{r1})E_a J_1(k_{r1}r) \sin \phi. \end{aligned} \quad (8)$$

To preserve the same acceleration field E_a as TM_{01} mode, the mixed mode requires a factor of two more power.

For the three modes we have $U_{11} = 2.405$, $U_{01} = U_{21} = 3.832$, $J_0(U_{01}) = -0.403$ and $J_1(U_{11}) = 0.519$. Coupling between the waveguide modes and Gauss-Laguerre modes, also known as TEM_{pl} modes [11], can be very efficient. When focused at waveguide entrance, power coupling from a radially polarized TEM_{01} mode to the TM_{01} mode reaches a maximum of 97% at $w_0/R = 0.56$. This is true also for coupling from a linearly polarized TEM_{01} mode to the mixed $TM_{01} + EH_{21}$ mode. Coupling from a TEM_{00} mode to the EH_{11} mode is 98% at $w_0/R = 0.64$, where w_0 is the Gaussian mode waist.

The on-axis intensity for a TEM_{00} mode falls as $I(z)/I(0) = 1/[1 + (z/Z_R)^2]$ away from the waist due to diffraction, and the on-axis longitudinal field for a TEM_{01} mode also falls as $E_z(z)/E_z(0) = 1/[1 + (z/Z_R)^2]$, where $Z_R = \pi w_0^2/\lambda_1$ is the Rayleigh length. Assuming a $TEM_{00}(TEM_{01})$ mode is coupled to a $EH_{11}(TM_{01})$ mode at the optimal condition, the effectiveness of guiding can be measured by taking the ratio of the relevant e-folding lengths of the waveguide mode to the free-space mode, yielding $L_{EH11}/L_{TEM00} = L_{TM01}/L_{TEM01} = 2R/\text{Re}(\Lambda)\lambda_1$. Despite the fact that the waveguide modes are leaky, optical guiding can be made quite effective to overcome diffraction for low order modes with sufficiently large R/λ_1 . In addition, for waveguide material with anomalous dispersion at certain wavelength, it is possible to have $\text{Re}(\Lambda) \ll 1$. For example, we have $\text{Re}(\Lambda) = 0.1$ for sapphire at $\lambda = 10.6\mu\text{m}$ with $\nu_2 = 0.67 + i0.04$ [17].

III. ACCELERATION IN VACUUM

According to Eq.(2), phase velocity of the TM_{01} mode is $v_p = \omega/\beta_{01} = c/(1 - 1/2\gamma_g^2)$, larger than the speed of light. We define an acceleration phase slippage length by

$$L_a = \frac{\lambda}{1/\gamma_g^2 + 1/\gamma^2}, \quad (9)$$

over which a relativistic electron with energy $W_0 = \gamma mc^2$, while being accelerated, slips π phase with respect to the fast acceleration wave. Energy gain of the electron on the axis is $\Delta W_a = eE_a \int_0^{L_a} \sin(\pi z/L_a) dz = eE_a L_a T_a$, where $T_a = 2/\pi$ is a transit factor. Here we have neglected the small waveguide attenuation over the distance L_a . In parallel, we may also define a deceleration slippage length over which the electron slips another π phase while losing energy amounted to $\Delta W_d = eE_a L_d T_d$. Average acceleration during a period of 2π phase slippage is then

$$G = \frac{\Delta W_a - \Delta W_d}{L_a + L_d} = \frac{eE_a T_a [1 - (L_d/L_a)(T_d/T_a)]}{1 + L_d/L_a}. \quad (10)$$

To have net acceleration, the ratio L_d/L_a should be made small. Two methods are proposed here for this purpose. The idea is to enhance phase slippage during the half period of deceleration, thus taking a shorter distance L_d .

The first method works on reducing the longitudinal velocity of an electron by introducing a static transverse magnetic field during deceleration. Without loss of generality, we assume the magnetic field is sinusoidal as in a wiggler $B_y = B_0 \cos(2\pi z/\lambda_w)$ with a period λ_w . Hence by definition L_d is determined through the condition

$$\left[\frac{1}{\gamma_g^2} + \frac{1}{\gamma^2} + \frac{a_w^2}{\gamma^2} \right] \frac{\pi L_d}{\lambda} - \frac{a_w^2 \lambda_w}{4\gamma^2 \lambda} \sin \left[\frac{4\pi L_d}{\lambda_w} \right] = \pi, \quad (11)$$

where $a_w = eB_0\lambda_w/2\pi\sqrt{2}mc$. If we set $\lambda_w = L_d$, then

$$L_d = \frac{\lambda}{1/\gamma_g^2 + 1/\gamma^2 + a_w^2/\gamma^2}, \quad (12)$$

here now $a_w = \sqrt[3]{Q_1 + \sqrt{Q_1^2 + Q_2^3}} + \sqrt[3]{Q_1 - \sqrt{Q_1^2 + Q_2^3}}$, with $Q_1 = eB_0\lambda\gamma^2/4\pi\sqrt{2}mc$ and $Q_2 = [1 + (\gamma/\gamma_g)^2]/3$. Due to longitudinal oscillation in electron orbit, $T_d/T_a = \frac{1}{2} \int_0^\pi \sin[\theta - \sin(4\theta)(1 - L_d/L_a)/4] d\theta$, which varies within the range $\{0.92 \leftrightarrow 1\}$ for L_d/L_a in the full range $\{0 \leftrightarrow 1\}$. The radiative energy loss per wiggler period is $\Delta W_r = (8\pi^2 mc^2/3)(r_e/\lambda_w)a_w^2\gamma^2$, where r_e is the classical radius of electron. The maximum possible energy that can be accelerated with this method is restricted by the condition $\Delta W_a > \Delta W_d + \Delta W_r$.

Table 1. Example in Vacuum with Magnetic Field.

λ [μm]	1	P_0 [TW]	50	E_a [GV/m]	1.8
R/λ	300	B_0 [T]	3	E_s [GV/m]	1.5
ν_2	1.5	a_w	22	I_s [TW/cm ²]	0.3
γ_g	492	$\Delta X_{max}/R$	0.37	G [GeV/m]	0.45
W_0 [GeV]	5	L_a [cm]	24	ΔW_a [MeV]	283
β_t [cm]	46	L_d [cm]	11	ΔW_d [MeV]	124
L_{attn} [m]	18	T_d/T_a	0.96	ΔW_r [keV]	16

Transverse EM force on a relativistic electron does not vanish to order of $1/\gamma^2$ in a waveguide mode or when the index of refraction differs from unity. For the TM_{01} mode with magnetic field $H_\phi = (1 + 1/2\gamma_g^2 + \delta\nu_1)E_r/Z_0$, transverse force is $F_r = -e(1/2\gamma^2 - 1/2\gamma_g^2 - \delta\nu_1)E_r$, which may lead to either focusing or defocusing depending on acceleration phase ϕ_a . Correspondingly, the beta function for an electron near the axis in vacuum is

$$\beta_t = \gamma_g \lambda \sqrt{(\gamma mc^2/\pi e \lambda E_a \sin \phi_a)/[1 - (\gamma_g/\gamma)^2]}. \quad (13)$$

The condition for an electron to remain near the axis may not be satisfied in the deceleration section, where the maximum orbital offset in x-direction is $\Delta X_{max} = \sqrt{2}a_w\lambda_w/\pi\gamma$. To avoid strong nonlinear transverse EM force while an electron is deflected far off-axis by the wiggler field, the mixed mode $TM_{01} + EH_{21}$ may be used,

for it has zero transverse field along x-axis. An example is given in Table 1 for acceleration in vacuum with TM_{01} mode, where in calculating β_t we set $\sin \phi_a = 1$, and I_s is laser intensity on waveguide inner surface.

Instead of tempering electron orbit, the second method works on enhancing phase slippage by increasing phase velocity of the wave during deceleration. This can be done by introducing a plasma layer of thickness

$$L_d = \frac{\lambda}{1/\gamma_g^2 + 1/\gamma^2 + 1/\gamma_p^2}, \quad (14)$$

where $\gamma_p = \omega/\omega_p \gg 1$, $\omega_p = c\sqrt{4\pi r_e n_0}$ is the electron plasma frequency and n_0 is the plasma density. Here we have used for the plasma an index of refraction $\nu_1 = 1 - \omega_p^2/2\omega^2$. In this case, $T_d = T_a$. The dominant energy loss for an ultrarelativistic electron traversing a plasma is due to bremsstrahlung [18]. The rate of energy loss is given by $dW/dz = -W/L_R$, where L_R is the radiation length defined by $1/L_R = 4\alpha r_e^2 n_i Z(Z+1) \ln(233/Z^{1/3})$, n_i is density of ions with atomic number Z , and α is the fine structure constant. For Hydrogen plasma with density $n_i = n_0 = 10^{17}/\text{cm}^3$, L_R is as long as $4 \times 10^6 \text{m}$. Reflectance of laser power off a sharp interface between vacuum and an underdense plasma at normal incidence is also negligible, according to the Fresnel formula $R_p = (1 - \nu_1)^2/(1 + \nu_1)^2 = 1/16\gamma_p^4$. Two examples are given in Table 2 and Table 3 for highly relativistic electron satisfying $(\gamma/\gamma_g)^2 \gg 1$ and $(\gamma/\gamma_p)^2 \gg 1$.

Table 2. Example in Vacuum with Plasma Layer (I).

λ [μm]	1	P_0 [TW]	10	E_a [GV/m]	1.9
R/λ	200	n_0 [$10^{17}/\text{cm}^3$]	1.1	E_s [GV/m]	1.5
ν_2	1.5	L_{attn} [m]	5.3	G [GeV/m]	1
γ_g	328	L_a [cm]	10.8	ΔW_a [MeV]	127
γ_p	100	L_d [cm]	0.91	ΔW_d [MeV]	11

Table 3. Example in Vacuum with Plasma Layer (II).

λ [μm]	10.6	P_0 [TW]	5	E_a [GV/m]	2.0
R/λ	50	n_0 [$10^{17}/\text{cm}^3$]	0.11	E_s [GV/m]	0.8
$\text{Re}(\nu_2)$	0.67	L_{attn} [m]	18	G [GeV/m]	1
$\text{Im}(\nu_2)$	0.04	L_a [cm]	7.1	ΔW_a [MeV]	90
γ_p	30	L_d [cm]	0.84	ΔW_d [MeV]	11

The assumption of a sharp vacuum-plasma interface is not necessary. More rigorous treatment can be obtained with WKB method [19]. For underdense plasma, the only modification required is to replace the factor $\exp(i\beta_{01}z)$ by $\exp[i\int \beta_{01}(z)dz]$. Assuming a density profile $n_e = n_0 f_p(z)$, where $f_p(z) = 1/[1 + e^{-(z+L_d/2)/\delta}] - 1/[1 + e^{-(z-L_d/2)/\delta}]$, the phase advance for the mode can be

calculated by making use of the integral $\int_{-\infty}^{\infty} f_p(z) dz = 2\delta \ln[(1 + e^{L_d/2\delta})/(1 + e^{-L_d/2\delta})]$. It is seen here that Eq.(14) is accurate enough as long as $L_d/2\delta \gg 1$. In addition, the validity of WKB method requires $|d\nu_1/dz| \ll 2\pi\nu_1^2/\lambda$ [19], which gives $\delta \gg \lambda/16\pi\gamma_p^2$ for $df_p/dz|_{max} = 1/4\delta$, also easily satisfied. To implement plasma layers in sections of a waveguide, side access is preferred. This can be done by inserting sections of open iris-loaded waveguide [20] wherever necessary. The mode profile in such a waveguide is nearly identical to that of a capillary waveguide, allowing significant minimization of mode coupling loss due to waveguide interruption.

IV. ACCELERATION IN GASES

The phase velocity of the TM_{01} mode in a core of uniform neutral gases is $v_p = \omega/\beta_{01} = c/(1 - 1/2\gamma_g^2 + \delta\nu_1)$, corresponding to a phase slippage length $L_a = \lambda/|1/\gamma_g^2 + 1/\gamma^2 - 2\delta\nu_1|$. The phase matching condition for synchronous acceleration is obtained by making the denominator zero, thus $\delta\nu_1 = 1/2\gamma_g^2 + 1/2\gamma^2$. This condition suggests an alternative way to maintain phase matching as γ increases during acceleration: instead of varying $\delta\nu_1$ by adjusting gas pressure along the waveguide, γ_g may be changed by tapering waveguide radius. For highly relativistic electron satisfying $(\gamma/\gamma_g)^2 \gg 1$, a steady-state phase matching condition, $\delta\nu_1 = 1/2\gamma_g^2$, is approached. An example is given in Table 4 in this limit. The beta function is smaller by a factor of $\sqrt{2}$ than that given for vacuum by Eq.(13). The maximum acceleration gradient achievable is limited by various processes occurring in gases in the field of high power laser, such as gas breakdown. Here we assume the limit is set on the maximum field for the TM_{01} mode by $E_r^{max} \leq 10\text{GV/m}$.

Table 4. Example in Gases.

λ [μm]	10.6	$\text{Re}(\nu_2)$	0.67	E_a [GV/m]	0.21
R/λ	50	$\text{Im}(\nu_2)$	0.04	E_s [GV/m]	0.085
P_0 [GW]	56	$\delta\nu_1[10^{-5}]$	7.4	L_{attn} [m]	18

Among other schemes of synchronous acceleration, one uses a line focus from an axicon in gases [7,8]. In comparison, our approach requires less laser power and provides much longer acceleration distance for the same γ , $\delta\nu_1$ and E_a . Another scheme uses a dielectrically lined perfect conductor tube as a lossless waveguide for TM modes [9,10]. This scheme, first of all, belongs to the class of acceleration by evanescent field near a surface of dielectric [3] and its scaling, $E_a/E_s = \lambda/\pi R$, is unfavorable for laser acceleration. Furthermore, the assumption of perfect conductor is no longer valid for metals at wavelength as short as $\lambda = 10.6\mu\text{m}$, where circular metallic waveguide is known to be very lossy for TM modes [18].

V. CONCLUSIONS

I have introduced the concepts and techniques that are crucial for advancing the current development of laser acceleration into a new, more realistic stage. Of what being accomplished here, the most notable is the dramatic increase in single-stage acceleration distance comparing with other schemes [4–10] for direct-field acceleration. This breakthrough is achieved in two critical steps. First, the diffraction limit is overcome with the capillary waveguide. Second, the phase slippage limit is overcome in vacuum with two new mechanisms for energy transfer. The two steps in combination succeed in harnessing high power laser to achieve high gradient with durable structure. In comparison, other schemes are severely limited for having components being blasted by full laser power in close range, head-on [4–6]. Furthermore, the approach taken here has the following advantages: oversized structure dimension is favorable for more accelerated particles; electromagnetically open waveguide suppresses higher order modes in favor of better beam quality; efficient coupling between waveguide and free-space modes simplifies mode handling such as injection, transport and recycling, leading to better overall system efficiency; last but not least, all aforementioned features are achieved without sacrificing a virtue of practical importance: the simplicity. Beam dynamics and collective effects will be addressed in a forthcoming series of papers. This work was supported by the U.S. Department of Energy under contract No.DE-AC03-76SF00098.

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