

A Proposed Dielectric-Loaded Resonant Laser Accelerator

J. Rosenzweig, A. Murokh, and C. Pellegrini

Department of Physics, University of California, Los Angeles, 405 Hilgard Avenue, Los Angeles, California 90024
(Received 2 September 1994)

A laser-pumped accelerator utilizing a resonant, periodic, dielectric structure is proposed. The electromagnetic fields due to a side-injected laser beam impinging on this structure are calculated in a two-dimensional standing-wave approximation, leading to an estimated accelerating field in excess of 1 GV/m for accessible experimental parameters. The longitudinal dynamics of injected electrons in the device are discussed, as are the first- and second-order transverse focusing effects inherent to this structure. Similarities and differences between this scheme and conventional and other advanced accelerators are examined.

PACS numbers: 29.17.+w, 41.75.Ht, 42.62.-b, 42.82.Et

Advanced acceleration techniques [1–7] are an active area of basic research, motivated by a range of applications spanning from high energy physics to radiation generation and medicine. In particular, there have been numerous advances in the acceleration of electrons using the large electric fields available in high power laser pulses. As the electric field is naturally transverse in a laser pulse, in order to impart energy gain secularly to an electron, either the electric field must be rotated to have a longitudinal component, as in the laser wake-field accelerator [1], the inverse Smith-Purcell effect [2], and the inverse Cerenkov accelerator [3], or the electron trajectory must have a periodic transverse component, as in the inverse free-electron laser [4]. These schemes have a common drawback, however, in that they demand very high peak laser power to obtain large acceleration gradients.

More promising schemes employ resonant excitation of a structure to build up large field amplitudes. Resonant accelerators utilize smaller peak powers to excite the acceleration fields, and can be efficiently beam loaded by long pulse trains, leading to higher average output currents and efficiencies. Perhaps the best example of this sort of scheme which uses lasers is the plasma beat-wave accelerator [5], in which the accelerating fields are provided by resonantly exciting electrostatic plasma waves. Unfortunately, because of the presence of free plasma electrons and associated currents inside of the beam channel, the transverse and longitudinal fields in this scheme are not of as high quality for accelerating low energy spread and emittance beams as conventional accelerators [6]. Problems in maintaining beam quality are not unique to plasma-based schemes, but are also present in laser accelerators with asymmetric geometries such as gratings [2].

To remove these objections to present laser-based acceleration schemes, we should therefore look for a resonantly excited, symmetric, purely electromagnetic structure. Plasma waves do not easily provide such a structure, since plasmas must be driven nonlinearly to

remove all the electrons from the beam channel [6]. A laser-pumped resonant accelerator may be easily constructed, however, using simple periodic metal or dielectric structures. As dielectric systems are attractive from the point of view of lower losses at optical frequencies, we will examine a dielectric-based system in this Letter.

To begin the description of the structure, we must first consider the trapping of laser power in a cavity, which can be thought of as the gap in a Fabry-Pérot interferometer made of a pair of highly reflective dielectric or metallic mirrors. In a pulsed system, the power incident on this mirror pair is initially nearly entirely reflected, while in steady state, the cavity fills with energy to the level where equilibrium is established, nearly all power is transmitted by the mirror-cavity system. A large stored energy U , and thus large electromagnetic field amplitudes in this cavity system, can be obtained with low external power P by having a large quality factor $Q = \omega U/P$. The Q of the cavity system operating at angular frequency ω also controls its exponential filling time, $\tau_f = Q/\omega$. In general, it is desirable to have a large Q , and filling time, but the laser pulse length may be limited by laser damage effects. Recent studies by Du *et al.* of optical damage in SiO_2 [7] have shown that for ultrashort laser pulses, $\tau_l \leq 1$ ps, the damage threshold of laser energy per unit area U_l/A has been found to scale as τ_l^{-1} , while for longer pulse lengths the threshold and the pulse length dependence of the threshold become $\tau_l^{1/2}$. This implies that the peak allowable power per unit area is proportional to τ_l^{-2} for $\tau_l \leq 1$ ps, and $\tau_l^{-1/2}$ for $\tau_l \geq 10$ ps, scalings which strongly favor shorter pulse lengths and thus the choice of Q and operating frequency. It is also important to note that a resonant laser accelerator, while requiring peak power scaling as Q^{-1} , would suffer breakdown at the same field level as a nonresonant structure, as it is the field amplitude which dictates the material damage through avalanche ionization. For an example relative to our present discussion, at $\tau_l = 2.5$ ps, the limiting energy flux is approximately 3 J/cm^2 [7], with an associated peak field amplitude of 2.5 GV/m.

To create an accelerating electromagnetic field profile, in addition to resonantly storing the power, the field component in the direction of the electron motion must be modulated. This modulation must have a spatial wave number equal to the free-space wave number of the laser light $k_0 = \omega/c$, in order to create a traveling wave spatial harmonic component [8] of the field having phase velocity equal to the speed of light which can secularly accelerate a relativistic ($v = \beta c \approx c$) electron. There are many ways of introducing a phase modulation, including periodic boundaries (e.g., an undulating wall), and period changes in the medium. Here we discuss a variation on the latter scheme, the use of a dielectric mask with a permittivity which is a periodic function of z , having a period $d = \lambda_0 = 2\pi/k_0$, the free-space wavelength. Furthermore, for the purpose of a model calculation, let us assume the following form of the permittivity within the mask:

$$\varepsilon(z) = \bar{\varepsilon} + (\Delta\varepsilon/2) \cos(k_0 z), \quad (1)$$

where $\bar{\varepsilon} = (\varepsilon_{\min} + \varepsilon_{\max})/2$ and $\Delta\varepsilon = \varepsilon_{\max} - \varepsilon_{\min}$. The geometry of this system, which is taken to be effectively infinite in the x dimension, is shown schematically in Fig. 1, with metallic mirrors chosen to allow the boundary condition outside of the outer dielectric to be approximated as perfectly conducting. This solution to this problem can easily be generalized to include any periodically varying permittivity.

In this model, we can find the symmetric TM modes of the structure by assuming a separable form of the longitudinal electric field $E_z = E_0 Y(y) Z(z) e^{i\omega t}$, ignoring the x dependence which occurs on a scale much larger than λ_0 . The familiar wave solution obtained for the y dependence of E_z is, assuming a field null at the mirror-mask boundary,

$$Y(y) = \sin\{k_y [y - (b + a)]\}, \quad (2)$$

where the vertical wave number k_y is an as yet undetermined separation constant. The z dependence of E_z within the dielectric is then given by

$$\frac{d^2 Z}{dz^2} + \left[k_0^2 \left(\bar{\varepsilon} + \frac{\Delta\varepsilon}{2} \cos(k_0 z) \right) - k_y^2 \right] Z = 0, \quad (3)$$

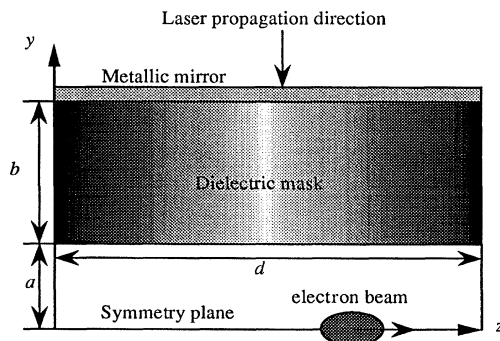


FIG. 1. Schematic representation of the top half of one period of a side-injected laser resonant acceleration structure using a modulated dielectric mask and a metallic mirror.

which can be recognized as the Mathieu equation. In order to maximize the first higher spatial (accelerating) harmonic, the solution to Eq. (3) of present interest must have two zeros in a period interval. This Mathieu function can be written approximately as [9]

$$Z(z) = \cos(k_0 z) - \Delta\varepsilon \left(\frac{\cos(2k_0 z)}{12} - \frac{1}{4} \right) + \Delta\varepsilon^2 \left[\frac{\cos(3k_0 z)}{384} - \frac{19 \cos(k_0 z)}{288} \right] + O(\Delta\varepsilon^3). \quad (4)$$

The amplitude of the component of this standing wave having phase velocity c (the forward wave $k_z = k_0$ spatial harmonic), relative to the other spatial harmonics, is maximized for $\Delta\varepsilon \ll 1$. From this point of view, it is best to choose $\Delta\varepsilon$ as small as possible, consistent with coupling of the incoming laser light to the TM accelerating mode, as this minimizes both the field amplitude and power loss associated with unuseful spatial harmonics. The details of coupling are beyond the scope of this Letter, but a simple estimate of optimum coupling can be obtained by requiring that the relative modulation of the phase of an incident plane wave due to differing optical path lengths between the regions of highest and lowest permittivity is maximized, $[\sqrt{\varepsilon_{\max}} - \sqrt{\varepsilon_{\min}}] k_0 b \approx \pi$.

Once the value of $\Delta\varepsilon$ is chosen based on these considerations, limiting the field to three significant spatial harmonics ($k_z = 0, \pm k_0$), the vertical wave number is then given by

$$k_y \approx k_0 \sqrt{\bar{\varepsilon} - 1 + \frac{5}{48} \Delta\varepsilon^2}. \quad (5)$$

If the thickness of the dielectric mask is also chosen to yield $k_y b = n\pi/2$, n odd, then $\partial E_z / \partial y = 0$ vanishes at the vacuum-dielectric interface. This allows for a smooth transition to the $k_z = \pm k_0$ spatial harmonic within the gap which, since they correspond to $k_y = 0$, are independent of y , and the acceleration due to the $k_z = \pm k_0$ spatial harmonic is independent of transverse position in the gap. The gap half-width a is constrained by the relation $k_0 a = m\pi$, m integer, which enforces mirror symmetry about the $y = 0$ plane of the fundamental ($k_y = k_0, k_z = 0$ in the free-space gap region) spatial harmonic. Inside of the gap, the largest component of the field is the $k_z = k_0$ standing wave whose forward component can secularly accelerate relativistic electrons with average force,

$$\bar{F}_z = -e\bar{E}_z \approx \sqrt{\frac{4\pi e^2 Q I}{c(1 + \Delta\varepsilon^2/8)}}, \quad (6)$$

where I is the incident laser intensity, the average is over a period of the structure, and we assume that the coupling of the laser to the cavity is unity.

An example list of parameters is given in Table I which describes the salient characteristics of this laser-based accelerator. The shortness of the laser pulse (which at 2.5 ps is chosen to give high breakdown threshold) compared to the electron transit time through the structure

TABLE I. Example parameters for a high gradient, resonant side-injected laser accelerator.

Laser wavelength (λ_0)	1.05 μm
Average permittivity ($\bar{\epsilon}$)	5.00
Permittivity contrast ($\Delta\epsilon$)	1.41
Dielectric mask thickness (b)	1.66 μm ($n = 15$)
Gap half-height (a)	1.575 μm ($m = 3$)
Laser illumination dimensions (l_x, l_z)	20 μm , 2 cm
Quality factor (Q)	500
Fill time (τ_f)	0.28 ps
Laser intensity (I)	8.4 GW/cm ²
Laser pulse length (τ_l)	2.5 ps
Average acceleration gradient ($e\bar{E}_z$)	1.0 GV/m

(67 ps) introduces an additional problem of requiring a correlation between the transverse and longitudinal laser intensity profile to ensure that the accelerating fields are present at the correct time in the structure. This can be accomplished through use of a dielectric or metallic “staircase” echelon to produce the necessary 45° angle in the laser intensity profile. Spatial filtering of the spurious transverse Fourier harmonics of the intensity profile generated by this component must be performed, however, as these harmonics affect the phase fronts of the laser light, and thus the acceleration process. The maintenance of coherence of the accelerating field is critical to this scheme, but this task is aided by the spatial coherence of a light source based on a laser oscillator, subsequent spatial filtering, and the rejection of temporal phase fluctuations by the resonant structure.

It should be noted that this type of accelerator is conceptually similar to a drift-tube linac [10], in that the component of the accelerating field which can secularly interact with a relativistic electron is created by periodically diminishing a zero-mode excitation. Although only the forward wave component of this standing wave imparts net acceleration to high energy electrons, the backward wave component provides second-order transverse focusing for these electrons, as discussed below.

The longitudinal dynamics of electrons in this structure are more interesting than in high gradient radio-frequency (rf) accelerators, because phase slippage is significant even for fairly relativistic electrons, due to the shortness of the wavelength. This slippage can be counteracted at low energies, by tapering the periodicity of the dielectric mask in z to satisfy $k_z(z) = 2\pi/d(z) = \omega/v(z)$, thus matching the phase velocity of the accelerating component of the field to the velocity of the electrons. If the field amplitude is also ramped positively in the transrelativistic injection region, then an initially nearly monochromatic, continuous beam of low energy electrons may be adiabatically captured into small regions of phase, allowing for very short microbunches with minimal energy spread. This type of capture and acceleration scheme is

in fact typical of proton linacs (such as an RFQ [10], or a drift-tube linac), where the transrelativistic region also has many spatial periodicity lengths.

In addition to longitudinal focusing, the transverse focusing associated with the capture section can be used to match the beam’s vertical phase space. This is due to two mechanisms arising from transverse electromagnetic forces, which for small vertical displacements can be written as $F_y \approx ey \partial E_z / \partial z$. There is a transverse focusing force first order in E_0 whenever there is a transient increase in this amplitude, and second-order focusing due to the electron encountering periodic field variations [11]. Generalizing the analysis in Ref. [11] to include the first-order focusing, and assuming transverse field variations only in y , we obtain an equation for the vertical motion of an ultrarelativistic electron,

$$y'' + \left(\frac{\gamma'}{\gamma}\right)\gamma' + \left[\frac{\gamma''}{\gamma} + \frac{1}{2}\left(\frac{\gamma'}{\gamma}\right)^2\right]y = 0. \quad (7)$$

Here the prime indicates the derivative with respect to distance along z , and we defined the average normalized energy gain $\gamma'(z) \equiv -e\bar{E}_z(z)/m_e c^2$. The term proportional to γ' corresponds to adiabatic damping of the motion, and the terms proportional to γ'' and $(\gamma')^2$ are due to the first- and second-order electromagnetic focusing, respectively.

As a practical example of this focusing, we take the parameters listed in Table I: $e\bar{E}_z = 1$ GeV/m ultimate average acceleration, rising linearly over 0.5 cm ($\gamma'' \approx \gamma'_{\text{max}}/L_r$). We note that the betatron wave numbers (k_β) associated with the first- and second-order focusing are approximately equal for the case $\gamma \approx 7$, with the combined focusing β function of 0.5 cm. This is a desirable condition, as when $\beta_f \approx L_r$ it is straightforward to match the beam’s transverse phase space to a focusing channel [12]. A matched beam envelope is indeed possible under conditions of constant acceleration, with rms beam size $\sigma_{y,\text{eq}} = (2)^{1/4} \sqrt{\epsilon_{yn}/\gamma'}$, where ϵ_{yn} is the normalized rms vertical emittance. To further illustrate this point, the envelope equation corresponding to the above example, with $\gamma_0 = 2$, $\gamma_f = 7$, and $\epsilon_{yn} = 10^{-10}$ m rad, is integrated, with the results shown in Fig. 2.

In this example, it is the electromagnetic focusing which allows the beam to propagate through the gap of the structure. The physical basis of this effect, however, is beneficial only in symmetric structures. In asymmetric systems such as the inverse Smith-Purcell accelerator, transverse electromagnetic forces will produce uncompensated deflections, not focusing effects. It is also apparent from the example that the ϵ_{yn} required for this scheme is very small, well below the transverse emittances injected into present electron rf linacs. This is a general feature of acceleration at optical wavelengths, as the beam size scales with the accelerating wavelength. Likewise, the pulse length and longitudinal emittance of an electron bunch adiabatically captured in an optical wave would be

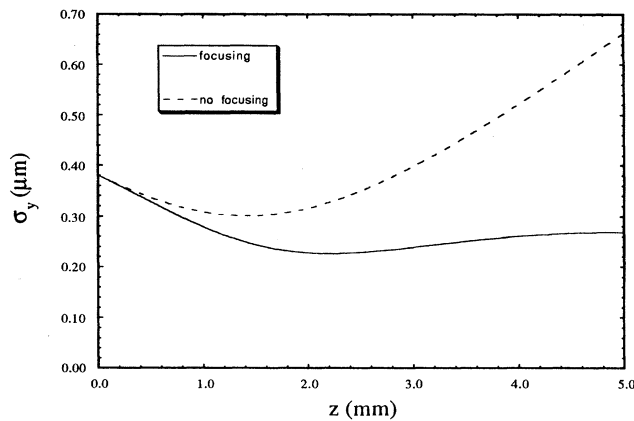


FIG. 2. Evolution of rms envelope (solid line) in a linearly rising field region ($L_r = 5$ mm) with all other parameters given in Table I. The dashed line gives the envelope in the absence of the focusing and accelerating electromagnetic field.

much smaller than that achieved in rf devices. This scaling can be viewed alternatively as a technical challenge, and as a desirable result of laser acceleration, since beams of such small phase space volume would undoubtedly lead to many advances in electron beam-based sciences.

In addition to independent particle dynamics, collective effects must be examined as well, most importantly the space-charge fields of the beam at low energy, and the beam-excited transverse modes in an accelerator, which can lead to an instability known as beam breakup (BBU) [13]. The geometry of the beam, which is much larger in the x than the y dimension, by itself will mitigate the space-charge effects in comparison to a round beam [10]. However, short wavelength structures inherently couple more strongly to the beam current due to the proximity of the walls of the structure to the beam axis. Since the structure under consideration is open horizontally and furthermore is overmoded, it only confines well the mode being externally pumped, and any beam excited transverse modes should radiate away quickly. Thus this structure should have a relatively high threshold current for BBU, which is important in allowing high levels of beam loading, and therefore high power efficiency.

The promise of the system we described in this Letter is to obtain a compact, laser excited, high gradient, inexpensive electron accelerator with attractive injection and transport properties. It is conceptually similar to present day accelerators, but with innovations in the structure and

coupling to adapt to much shorter wavelength operation. Most of the demands on the laser system, which in the example uses only $200 \mu\text{J}$ per pulse, can be met with commercial products. We note that, while our example assumes use of a $1 \mu\text{m}$ laser, initial proof-of-principle experiments should be performed at longer wavelengths, easing the structure manufacturing, as well as the stringent injected electron source requirements. However, there is no reason to believe that these structures are overly challenging to build; nanofabrication techniques developed in recent years, for advanced applications such as gradient index optics [14] and microcavity lasers [15], can also meet most of the requirements for construction of the dielectric masks.

The authors acknowledge useful and stimulating conversations with R. Byer, W. Gai, S. T. Ho, W. Kimura, and N. Kroll. This work was performed with partial support from U.S. Department of Energy Grants No. DE-FG03-90ER40796 and No. DE-FG03-92ER40693, and the Alfred P. Sloan Foundation Grant No. BR-3225.

-
- [1] P. Sprangle, E. Esarey, A. Ting, and G. Joyce, *Appl. Phys. Lett.* **53**, 2146 (1988).
 - [2] R. Palmer *et al.*, in *Laser Acceleration of Particles*, edited by C. Joshi and T. Katsouleas, AIP Conf. Proc. 130 (AIP, New York, 1985), pp. 234–252.
 - [3] J. R. Fontana and R. H. Pantell, *J. Appl. Phys.* **54**, 4285 (1983).
 - [4] R. B. Palmer, *J. Appl. Phys.* **43**, 3014 (1972).
 - [5] C. E. Clayton *et al.*, *Phys. Rev. Lett.* **70**, 37 (1993).
 - [6] J. B. Rosenzweig *et al.*, *Phys. Rev. A* **44**, R6189 (1991).
 - [7] D. Du *et al.*, *Appl. Phys. Lett.* **64**, 3073 (1994).
 - [8] G. Loew and R. Neal, in *Linear Accelerators*, edited by Pierre M. Lapostolle and Albert L. Septier (North-Holland, Amsterdam, 1969).
 - [9] M. Abramowitz and I. Stegun, *Handbook of Mathematical Functions* (Dover, New York, 1964).
 - [10] P. Lapostolle, *Proton Linear Accelerators: A Theoretical and Historical Introduction* (LANL Pub. LA-11601-MS, Los Alamos, 1989).
 - [11] J. Rosenzweig and L. Serafini, *Phys. Rev. E* **49**, 1599 (1994).
 - [12] N. Barov and J. Rosenzweig, *Phys. Rev. E* **49**, 4407 (1994).
 - [13] Y. Y. Lau, *Phys. Rev. Lett.* **63**, 1141 (1989).
 - [14] Proceedings of GRIN '92 [*Appl. Opt.* **33**, No. 16 (1994)].
 - [15] Y. Yamamoto and R. Slusher, *Phys. Today* **46**, No. 6, 66 (1993).

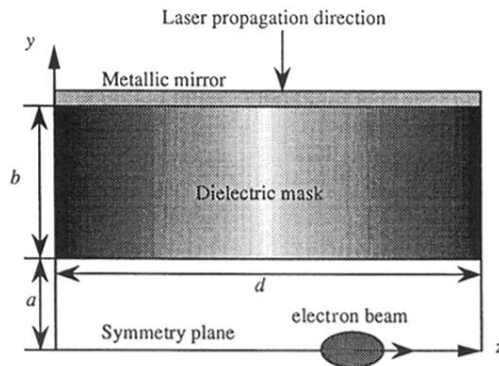


FIG. 1. Schematic representation of the top half of one period of a side-injected laser resonant acceleration structure using a modulated dielectric mask and a metallic mirror.