

Electron acceleration in vacuum by ultrashort and tightly focused radially polarized laser pulses

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Abstract. Exact closed-form solutions to Maxwell's equations are used to investigate electron acceleration driven by radially polarized laser beams in the nonparaxial and ultrashort pulse regime. Besides allowing for higher energy gains, such beams could generate synchronized counterpropagating electron bunches.

1 Introduction

The advent of ultra-intense laser facilities has led to exciting possibilities in the development of a new generation of compact laser-driven electron accelerators. Among the proposed laser acceleration schemes, the use of ultra-intense radially polarized laser beams in vacuum (termed *direct acceleration*) is very promising, as it takes advantage of the strong longitudinal electric field at beam center to accelerate electrons along the optical axis [1]. Numerical simulations have shown that collimated attosecond electron pulses could be produced by this acceleration scheme [2, 3].

Recent studies on direct acceleration have shown that reducing the pulse duration and beam waist size generally increases the maximum energy gain available [4, 5]. However, these analyses were carried under the paraxial and slowly varying envelope approximations. These approximations lose their validity as the beam waist size becomes comparable with the laser wavelength and the pulse duration approaches the single-cycle limit, conditions that are now often encountered in experiments. We propose a simple method to investigate direct acceleration in the nonparaxial and ultrashort pulse regime, and show that it offers the possibility of higher energy gains. We also highlight a peculiar feature of the acceleration dynamics under nonparaxial focusing conditions, namely the coexistence of forward and backward acceleration. This could offer a solution to the production of synchronized electron pulses required in some pump-probe experiments.

2 Exact solution for a nonparaxial and ultrashort TM_{01} pulsed beam

Ultrashort and tightly focused pulsed beams must be modeled as exact solutions to Maxwell's equations. A simple and complete strategy to obtain exact closed-form solutions for the electromagnetic fields of such beams was recently presented by April [6]. For a TM_{01} pulse, which corresponds to the lowest-order radially polarized laser beam, the field components are described by [6, 7]:

$$E_r(\mathbf{r}, t) = \text{Re} \left\{ \frac{3E_0 \sin 2\tilde{\theta}}{2\tilde{R}} \left(\frac{G_-^{(0)}}{\tilde{R}^2} - \frac{G_+^{(1)}}{c\tilde{R}} + \frac{G_-^{(2)}}{3c^2} \right) \right\}, \quad (1)$$

$$E_z(\mathbf{r}, t) = \text{Re} \left\{ \frac{E_0}{\tilde{R}} \left[\frac{(3 \cos^2 \tilde{\theta} - 1)}{\tilde{R}} \left(\frac{G_-^{(0)}}{\tilde{R}} - \frac{G_+^{(1)}}{c} \right) - \frac{\sin^2 \tilde{\theta}}{c^2} G_-^{(2)} \right] \right\}, \quad (2)$$

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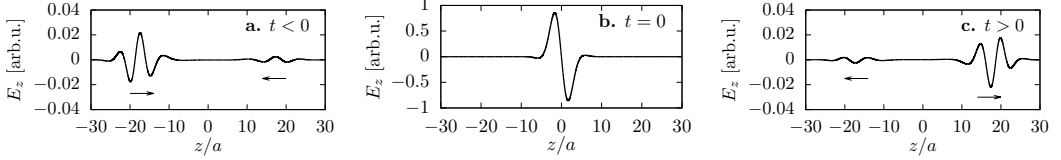


Fig. 1. Longitudinal on-axis electric field of a TM_{01} pulse with $k_0a = 1$ and $s = 10$.

$$H_\phi(\mathbf{r}, t) = \text{Re} \left\{ \frac{E_0 \sin \tilde{\theta}}{\eta_0 \tilde{R}} \left(\frac{G_-^{(1)}}{c \tilde{R}} - \frac{G_+^{(2)}}{c^2} \right) \right\}. \quad (3)$$

Here E_0 is an amplitude parameter, $\tilde{R} = [r^2 + (z + ja)^2]^{1/2}$, $\cos \tilde{\theta} = (z + ja)/\tilde{R}$, $G_\pm^{(n)} = \partial_t^n [f(\tilde{t}_\pm) \pm f(\tilde{t}_-)]$, $f(t) = e^{-j\phi_0} (1 - j\omega_0 t/s)^{-(s+1)}$, and $\tilde{t}_\pm = t \pm \tilde{R}/c + ja/c$. The function $f(t)$ is the inverse Fourier transform of the Poisson-like frequency spectrum of the pulse, in which $\omega_0 = ck_0$ is the frequency of maximum amplitude and ϕ_0 is a constant phase [8]. The parameter a , called the confocal parameter, is monotonically related to the beam waist size and characterizes the beam's degree of paraxiality: $k_0a \sim 1$ for tight focusing conditions, while $k_0a \gg 1$ for paraxial beams. The pulse duration T , which may be defined as twice the root-mean-square width of $|E_z|^2$, increases monotonically with s . In the limit $k_0a \gg 1$ and $s \gg 1$, Eqs. (1)–(3) reduce to the familiar paraxial TM_{01} Gaussian pulse [9].

The TM_{01} pulsed beam described above may be produced by focusing a collimated radially polarized input beam with a high aperture parabolic mirror. Its field distribution consists of two counter-propagating pulse components, as shown in Fig. 1 [10].

3 On-axis acceleration in the nonparaxial and ultrashort pulse regime

Direct acceleration is simulated by integrating the conventional Lorentz force equation for an electron initially at rest at position z_0 on the optical axis and outside the laser pulse. Since E_r and H_ϕ vanish at $r = 0$, the particle is accelerated by E_z along the optical axis.

Figure 2 illustrates the variation of the maximum energy gain available ΔW_{\max} (after optimizing for z_0 and ϕ_0) with the laser peak power P_{peak} for different combinations of k_0a and s . Figure 2a, in which ΔW_{\max} is expressed as a fraction of the theoretical energy gain limit ΔW_{lim} [9], shows that for constant values of s , the threshold power above which significant acceleration occurs is greatly reduced as k_0a decreases, i.e., as the focus is made tighter. According to Fig. 2b, MeV energy gains may be reached under tight focusing conditions with laser peak powers as low as 15 gigawatts. In contrast, a peak power about 10^3 times greater is required to reach the same energy with paraxial pulses. At high peak power, Fig. 2a shows that shorter pulses yield a more efficient acceleration, with a ratio $\Delta W_{\max}/\Delta W_{\text{lim}}$ reaching 80% for single-cycle ($s = 1$) pulses. Additional details about those results can be found in [7].

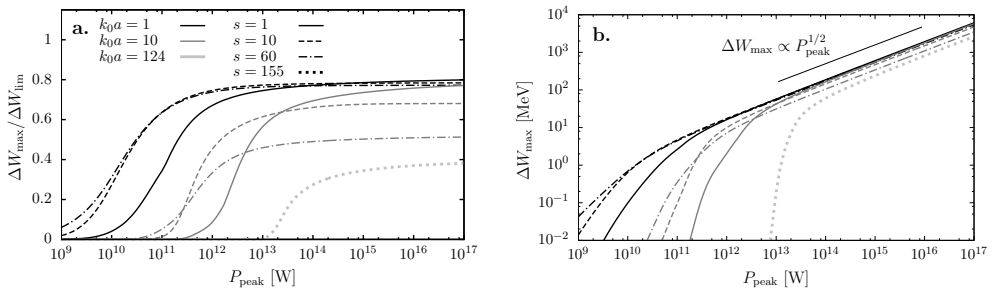


Fig. 2. Maximum (a) normalized and (b) absolute energy gain of an electron initially at rest versus the laser pulse peak power for different values of k_0a and s . The curve with $\{k_0a = 124, s = 155\}$ corresponds to the limit of the paraxial regime investigated in [4]. Figure taken from [7].

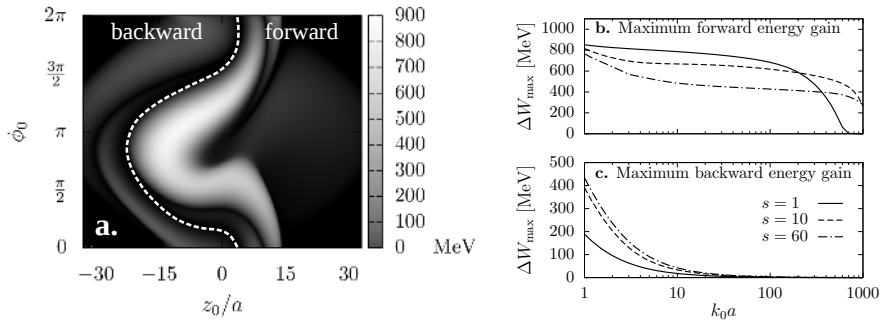


Fig. 3. (a) Energy gain of an electron initially at rest versus z_0 and ϕ_0 for a laser pulse with $k_0a = 1$, $s = 1$. The dashed curve delimits the regions of forward and backward acceleration. (b)–(c) Maximum energy gain for electrons accelerated forward and backward versus k_0a . In all figures, $P_{\text{peak}} = 2 \times 10^{15}$ W.

In the highly nonparaxial regime ($k_0a \sim 1$), a closer look at the dynamics in the (z_0, ϕ_0) parameter space reveals the existence of two different types of acceleration (see Fig. 3a). In the first type, the electron is accelerated in the positive z direction (forward acceleration), and may reach a high energy gain if its motion is synchronized with a negative half-cycle of the forward-propagating component of the beam. In the second type, the electron is accelerated in the negative z direction (backward acceleration), and may similarly experience subcycle acceleration from the backward-propagating component of the beam. The maximum energy gain available from forward and backward acceleration is illustrated in Figs. 3b–c. A significant backward acceleration is only observed under tight focusing conditions ($k_0a < 10$), since the amplitude of the backward-propagating component of the laser beam rapidly decreases as k_0a increases.

4 Conclusion

We have highlighted the importance of going beyond the paraxial and slowly varying envelope approximations in the analysis of electron acceleration in vacuum by radially polarized laser beams. It was shown that the acceleration threshold power may be greatly reduced under tight focusing conditions, which demonstrates that direct acceleration is much more accessible to the current laser technology than previously expected. Moreover, our results hints that high-aperture focusing optics could be used to generate synchronized counterpropagating electron bunches. The proposed acceleration scheme could therefore find applications in the context of pump-probe experiments.

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