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A "slingshot" laser-driven acceleration mechanism of plasma electrons

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ABSTRACT

We briefly report on the recently proposed Fiore et al. [1] and Fiore and De Nicola [2] electron acceleration mechanism named "slingshot effect": under suitable conditions the impact of an ultra-short and ultra-intense laser pulse against the surface of a low-density plasma is expected to cause the expulsion of a bunch of superficial electrons with high energy in the direction opposite to that of the pulse propagation; this is due to the interplay of the huge ponderomotive force, huge longitudinal field arising from charge separation, and the finite size of the laser spot.

1. Introduction

Today ultra-intense laser–plasma interactions allow extremely compact acceleration mechanisms of charged particles to relativistic regimes, with numerous and extremely important potential applications in nuclear medicine (cancer therapy and diagnostics), research (particle physics, inertial nuclear fusion, optycs, materials science, structural biology, etc.), food sterilization, transmutation of nuclear wastes, etc. A prominent mechanism for electrons is the *wake-field acceleration* (WFA) [3]: electrons are accelerated "surfing" a plasma wake wave driven by a short laser or charged particle beam within a low-density plasma sample (or matter to be locally completely ionized into a plasma by the beam, more precisely a supersonic gas jet), and are expelled just after the exit of the beam out of the plasma, behind and *in the same direction as the beam* (forward expulsion). WFA has proved to be particularly effective since 2004 in the so-called *bubble* (or *blowout*) regime; it can produce electron bunches of very good collimation, small energy spread and energies of up to hundreds of MeVs [4–6] or more recently even GeVs [7,8].

In Refs. [1,2] it has been claimed that the impact of a very short and intense laser pulse in the form of a pancake normally onto the surface of a low-density plasma may induce also the acceleration and expulsion of electrons *backwards* (*slingshot effect*), see Fig. 1. A bunch of plasma electrons (in a thin layer just beyond the vacuum–plasma interface) first are displaced forward with respect to the ions by the positive ponderomotive force $F_p \equiv \langle -e(\frac{\mathbf{v}}{c}\times\mathbf{B})^z\rangle$ generated by the pulse (here $\langle\rangle$ is the average over a period of the laser carrier wave, $\mathbf{E},\mathbf{B}$ are the electric and magnetic fields, $\mathbf{v}$ is the electron velocity, and $\hat{\mathbf{z}}$ is the direction of propagation of the laser pulse; recall that F_p is positive, negative when the modulating amplitude ϵ_s of the pulse respectively increases, decreases), then are pulled back by the electric force $-eE^z$ due to this charge displacement. If the electron density $\widetilde{n_0}$ is tuned in the range where the plasma oscillation period T_H is about twice the pulse duration τ, then these electrons invert their motion when they are reached by the maximum of ϵ_s, so that the negative part of F_p adds to $-eE^z$ in accelerating them backwards; equivalently, the total work $W \equiv \int_0^\tau dt\, F_p v^z$ done by the ponderomotive force is maximal.[1] Their expulsion energy (out of the bulk) will be enough to escape to $z\to-\infty$ if the laser spot radius R is suitably tuned.

The very short pulse duration τ and expulsion time t_e, as well as huge nonlinearities, make approximation schemes based on Fourier analysis and related methods inconvenient. But recourse to full kinetic theory is not necessary: we show [9,2] that in the relevant space–time region a Magneto Hydro Dynamic (MHD) description of the impact is self-consistent, simple and predictive. The set-up is as follows. We describe the plasma as consisting of a static background of ions and a fully relativistic, collisionless fluid of electrons, with the system "plasma + electromagnetic field" fulfilling the Lorentz–Maxwell and the continuity Partial Differential Equations (PDE). For brevity, below we refer to the

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[1] Whereas $F_p v^z$ oscillates many times about 0, and $W \simeq 0$, if $\tau \gg T_H$ – the standard experimental situation until a couple of decades ago.

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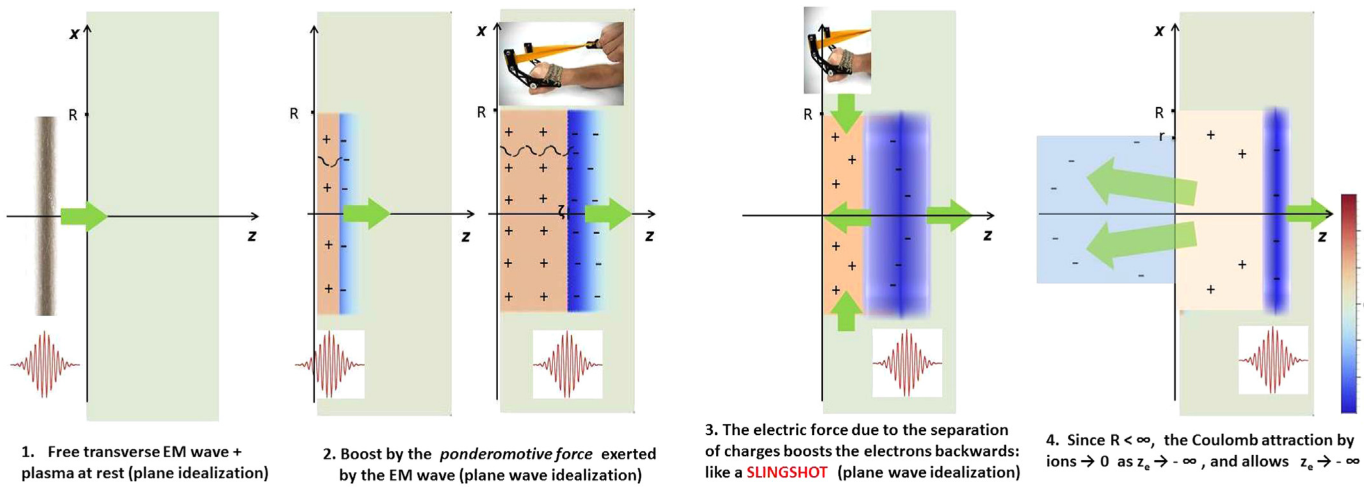

Fig. 1. Schematic stages of the slingshot effect.

electrons' fluid element initially located at $\mathbf{X} \equiv (X, Y, Z)$ as to the "$\mathbf{X}$ electrons", and to the fluid elements with arbitrary X, Y and specified Z as the "Z electrons". We denote: as $\mathbf{x}_e(t, \mathbf{X})$ the position at time t of the $\mathbf{X}$ electrons, and for each fixed t as $\mathbf{X}_e(t, \mathbf{x})$ the inverse of $\mathbf{x}_e(t, \mathbf{X})$ $[\mathbf{x} \equiv (x, y, z)]$; as c the velocity of light; as m and as $n, \mathbf{v}, \mathbf{p}$ the electrons' mass and Eulerian density, velocity, momentum. $\boldsymbol{\beta} \equiv \mathbf{v}/c$, $\mathbf{u} \equiv \mathbf{p}/mc = \boldsymbol{\beta}/\sqrt{1-\boldsymbol{\beta}^2}$, $\gamma \equiv 1/\sqrt{1-\boldsymbol{\beta}^2} = \sqrt{1+\mathbf{u}^2}$ are dimensionless. We assume that the plasma is initially neutral, unmagnetized and at rest with electron (and proton) density $\widetilde{n_0}(z)$ depending only on z and equal to zero in the region $z < 0$. We schematize the laser pulse as a free transverse EM plane travelling-wave multiplied by a cylindrically symmetric "cutoff" function, e.g.

$$\mathbf{E}^{\perp}(t, \mathbf{x}) = \boldsymbol{\epsilon}^{\perp}(ct - z)\theta(R - \rho), \quad \mathbf{B}^{\perp} = \hat{\mathbf{z}} \times \mathbf{E}^{\perp} \tag{1}$$

where $\rho \equiv \sqrt{x^2 + y^2} \le R$, θ is the Heaviside step function, and the "pump" function $\boldsymbol{\epsilon}^{\perp}(\xi)$ vanishes outside some finite interval $0 < \xi < l$. Then, to simplify the problem,

1. We first study the $R = \infty$ (i.e. *plane-symmetric*) version, carefully choosing unknowns and independent variables (Section 2.1). For sufficiently small densities and short times we can reduce the PDE's to a collection of decoupled *systems of two first order nonlinear ODE in Hamiltonian form*, which we solve numerically.
2. We determine (Section 2.2): $R < \infty$, $r > 0$ so that the plane version gives small errors for the surface electrons with $\rho \le r \le R$; the corresponding final energy, spectrum, etc. of the expelled electrons. For definiteness, we consider the $\widetilde{n_0}(z)$ of Fig. 2.

We specialize our predictions to virtual experiments at the FLAME facility (LNF, Frascati). We invite for simulations (PIC, etc.) and experiments testing them.

2. The model

2.1. Plane wave idealization

Our plane wave Ansatz reads: $A^{\mu}, \mathbf{u}, n - \widetilde{n_0}(z)$ depend only on z, t and vanish if $ct \le z$; $\Delta\mathbf{x}_e \equiv \mathbf{x}_e - \mathbf{X}$ depends only on Z, t and vanishes if $ct \le Z$. Then: $\mathbf{B} = \mathbf{B}^{\perp} = \hat{\mathbf{z}}\partial_z \wedge \mathbf{A}^{\perp}$, $c\mathbf{E}^{\perp} = -\partial_t \mathbf{A}^{\perp}$; the transverse component of the Lorentz equation implies $\mathbf{u}^{\perp} = e\mathbf{A}^{\perp}/mc^2$; due to

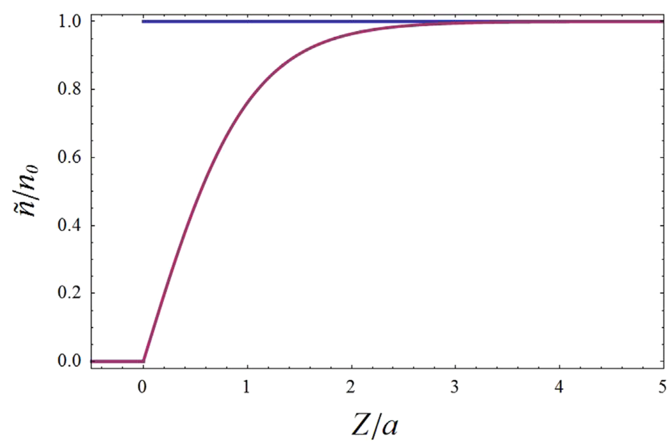

Fig. 2. The normalized $\widetilde{n}_0$ adopted here: step-shaped (blue) and continuous $\widetilde{n}_0(Z) = n_0\theta(Z)\tanh(Z/a)$, $a = 20\ \mu\mathrm{m}$ (purple); they respectively model the initial electron densities at the vacuum interfaces of an aerogel and of a gas jet (just outside the nozzle). (For interpretation of the references to color in this figure caption, the reader is referred to the web version of this paper.)

charge separation $E^z \ne 0$: by the Maxwell equations it is related to the longitudinal motion through

$$E^z(t, z) = 4\pi e\left\{\widetilde{N}(z) - \widetilde{N}[Z_e(t, z)]\right\}, \quad \widetilde{N}(Z) \equiv \int_0^Z d\eta\, \widetilde{n}_0(\eta) \tag{2}$$

what yields a conservative force on the electrons. For sufficiently small densities and short times the laser pulse is not significantly affected by the interaction with the plasma (the validity of this approximation is checked a posteriori [2]), and we can identify $\mathbf{A}^{\perp}(t, z) = \boldsymbol{\alpha}(\xi)$, $\xi \equiv ct - z$, where $\boldsymbol{\alpha}$ is the transverse vector potential of the "pump" free laser pulse. Hence also $\mathbf{u}^{\perp}(t, z) = e\boldsymbol{\alpha}(\xi)/mc^2$ is explicitly determined. For each fixed Z, the unknown $z_e(t, Z)$ appears in place of z in the equations of motion of the Z-electrons. But, as no particle can reach the speed of light, the map $t \mapsto \xi \equiv ct - z_e(t, Z)$ is strictly increasing, and we can use [2,9] (ξ, Z) instead of (t, Z) as independent variables. It is also convenient to use the "electron s-factor" $s \equiv \gamma - u^z$ instead of u^z as an unknown, because it is *insensitive* to rapid oscillations of $\boldsymbol{\alpha}$, and $\gamma, \mathbf{u}, \boldsymbol{\beta}$ are *rational* functions of $\mathbf{u}^{\perp}, s$:

$$\gamma = \frac{1 + \mathbf{u}^{\perp 2} + s^2}{2s}, \quad u^z = \frac{1 + \mathbf{u}^{\perp 2} - s^2}{2s}, \quad \boldsymbol{\beta} = \frac{\mathbf{u}}{\gamma}. \tag{3}$$