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Debye-Length Discrimination of Nonlinear Laser Forces Acting on Electrons in Tenuous Plasmas

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Experimental results are presented which reveal the important role played by nonlinear radiation or ponderomotive forces in laser plasma interactions. The existence of these nonlinear forces are clearly revealed by the properties of tenuous helium plasmas generated by focusing of intense laser beams.

One of us (H. H.) has drawn attention to the potentially important role[1,2] that the well-known "ponderomotive force" could play in laser-plasma interactions. However, because other, more established, mechanisms could also account for the properties of laser-produced plasmas observed during recent years, some uncertainty has always existed regarding the true role of nonlinear radiation forces[1,2] in such plasmas.

The important role of the nonlinear radiation forces[1,2] in modifying the density profile of dense, laser-produced plasmas was recently reported by Enright and Richardson[3] and by Azechi *et al.*[4] in this journal. These observations have had a critical bearing on laser fusion studies based on CO_2-laser excitation due to the fact that the nonlinear radiation forces[1,2] modify the plasma density profile in such a manner that the laser energy can be deposited much nearer to the critical-density surface than was previously thought possible.

In an effort to further the understanding of the properties of nonlinear radiation forces,[1,2] we have initiated a research program aimed at studying the properties of these important forces[5,6] under conditions where other effects, considered in the past to be the dominant processes in laser-plasma interactions, cannot possibly be effective. We believe that tenuous plasmas irradiated with intense laser radiation provides the ideal environment for the detailed study of the nonlinear radiation forces.[1,2] In this Letter we present experimental results which clearly demonstrate the manner in which the importance of the nonlinear radiation forces[1,2] are revealed as the density of the tenuous plasma is decreased.

The experimental arrangement is shown in Fig. 1. A 4 GW (0.1 J in 25×10^{-12} sec) beam of Gaussian intensity profile, generated by a Quentron Model 100, high-power, shoft rod Nd:YAlG:glass (yttrium aluminum garnet) laser system, was focused, using a SORO $F1.5$ aspheric lens, inside a chamber filled with helium at pressures in the range 10^{-2} to 10^{-5} Torr after initial evacuation to less than 5×10^{-7} Torr.

A monodirectional, retarding-field electron energy analyzer placed along the direction of the electric field vector allowed any electrons emit-

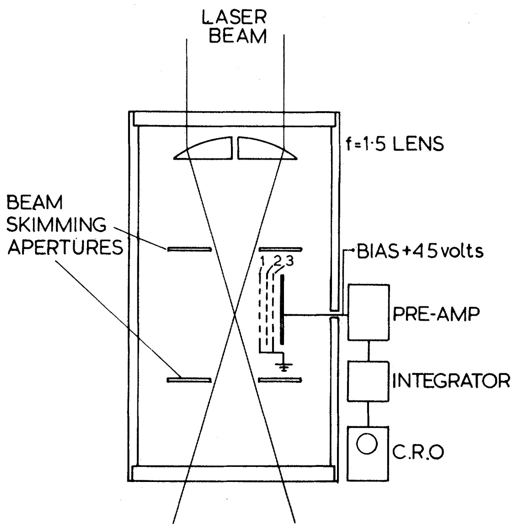

FIG. 1. Experiment with laser focus in vacuum chamber containing helium of pressures 10^{-6} to 10^{-3} Torr. 1, 2, and 3 are grids in front of the probe.

ted from the focal region of the lens to be collected over a solid angle of 1.5 sr. Grids 1 and 3 (Fig. 1) were grounded to produce a field-free region at the focus and Grid 2 acted as a retarding grid which was negatively biased to prevent electrons with energies less than the grid potential from reaching the collector. The collector itself was positively biased at + 45 V to prevent secondary electron losses. Beam skimming apertures were used to shield the collector from stray electrons originating at the lens and window surfaces and signal levels from these sources and also from the residual gas in the chamber were then found to be negligible.

The maximum charge which could be detected was ~ 2×10^{-15} C which is equivalent to ~ 10^5 electrons being generated in the focal region. A consequence of this high detection threshold is that electrons generated in the highest-intensity region of the beam will not be detected. For example, at a pressure of 10^{-4} Torr the volume of fully ionized gas must exceed 3×10^{-8} cm before any electrons can be detected and beam intensity distribution measurements show that the local beam intensity will have dropped to a factor of 3 below its peak value in the focal plane before the ionized volume exceeds this limit.[6]

We take the nonlinear radiation force[1,2] to be of the form[7]

$$\vec{f}_{NL} = -\frac{1}{8\pi}\nabla(E^2+H^2) = \frac{1-\tilde{n}^2}{8\pi}\nabla E^2 \quad (1)$$

within the radial intensity gradient of our focused laser beam where helium atoms are ionized in a period of $<10^{-11}$ sec with beam intensities exceeding 10^{14} W cm^{-2} at 1.06 μm.[6] In the case of free electrons under the influence of the nonlinear radiation force acting along the radial intensity gradient of our focused laser beam, the average kinetic energy $\epsilon_{kin,e}^{osc}$, of the maximum oscillation of the electrons, given by[7]

$$\epsilon_{kin,e}^{osc} = \tfrac{1}{2}\epsilon_e^{osc} = \frac{E^2}{16\pi n_{ec}|\tilde{n}|}, \quad (2)$$

is transferred into translational energy as the electron is accelerated out of the beam. Here n_{ec} is the critical electron density. It follows that the relation

$$\epsilon_e^{trans} = \epsilon_{kin,e}^{osc} \quad (3)$$

will be valid as long as free electrons are being accelerated by the nonlinear radiation forces. In our experiments, this requirement is satisfied as long as the "Debye length" λ_D^* corresponding to a "temperature" equal to this energy is larger than the diameter d, of our focus spot, i.e.,

$$\lambda_D^* = 734\left|\frac{2}{3}\frac{\epsilon_{kin,e}^{osc}(\text{eV})}{n_e(\text{cm}^{-3})}\right|^{1/2}\text{cm, for } \lambda_D^* > d, \quad (4)$$

where n_e is the electron density. When λ_D^* becomes less than d, electron-ion coupling occurs resulting in some of the electron energy being transferred to ion motion and only a few electrons of the Debye sheath can be accelerated as an electron could and be detected by the probe. On the other hand, ions of charge z will be accelerated via the electrostatic coupling to the electrons and will attain a maximum translation energy given by

$$\epsilon_i^{trans} = z\epsilon_{kin,e}^{osc}. \quad (5)$$

Assuming an ionized helium plasma we can relate the electron density n_e to a gaseous pressure p and relating $\epsilon_{kin,e}^{osc}$ to the laser power P and the beam diameter d, we obtain a relation for the transition at pressures $p=p^*$ from pure electron acceleration ($p<p^*$) to plasma acceleration ($p>p^*$) at $\lambda_D^*=d$ [Eq. (4)]. I.e.,

$$p^* = 1.37\times10^{-24}zP/d^4, \quad (6)$$

where P is the laser power in watts, d is in centimeters squared, and p^* in Torr. Figure 2 shows the detector signals for the situation $p>p^*$ where the upper trace is the ion signal and the lower trace the electron signal. No ion signal is obtained for the case with $p<p^*$.

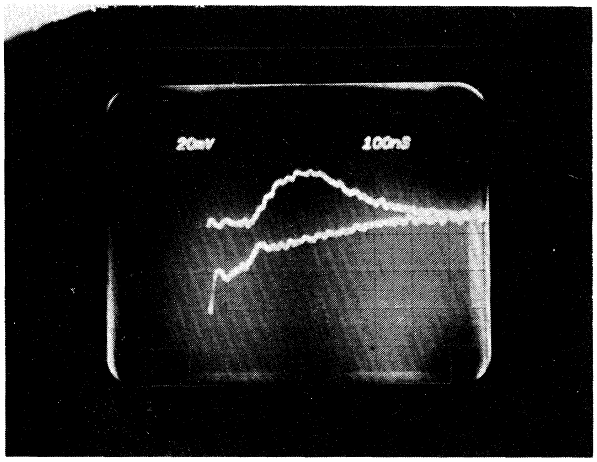

FIG. 2. Probe signal of electron and ion pulses at a pressure of 8×10^{-3} Torr.

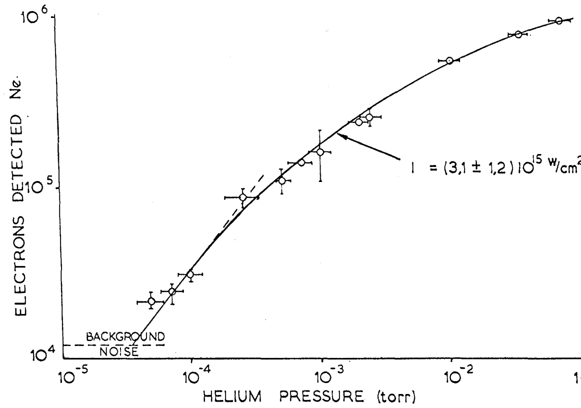

FIG. 3. Number of emitted electrons with energies above 50 eV from laser-induced gas breakdown in He of varying pressure P. Laser intensity I equals $(1.1 \pm 0.4)\times10^{15}$ W/cm^2.

To obtain more details of the processes taking place in our experiments, we measured electron yields as a function of gas pressure over the range 10^{-1} to 10^{-5} Torr and also as a function of retarding grid potential. From the above consideration it is clear that maximum electron energy will be generated when $p < p^*$. Under these conditions the energy of the detected electron will be identical to its energy as it emerges from the focus because the mean free path of the electron in this pressure range can be several orders of magnitude larger than the distance to the probes. The maximum electron energy measured at a pressure of 10^{-4} Torr was 90 ± 5 eV. Since the plasma generated in the focus under these conditions has a Debye length $\lambda_D{}^*$ larger than the focus, we can assume that the nonlinear radiation force expels the electrons from regions of maximum intensity within the laser beam converting the average kinetic energy of quivering motion to directed kinetic energy. From Eqs. (2) and (3) with $|\tilde{n}|\sim1$ (very dilute plasma), $d = 13\pm2$ μm (measured in a separate experiment[6,8]) and an effective peak intensity of 10^{15} W/cm^2 we obtain $\epsilon_e{}^{\text{trans}} = 104\pm32$ eV, which agrees well with the experimentally measured value within the accuracy of the measurement.

The number of electrons, N_e, collected with energies greater than 50 eV is plotted as a function of the helium pressure in Fig. 3. This curve would be expected to have a linear slope for pressures where $\lambda_D{}^* \gg d$ and it is seen that this is the case for a pressure of $(1.6\pm0.5)\times10^{-4}$ Torr. A similar result has been obtained in experiments in nitrogen gas.[9] Detector limitations prevented a similar curve being obtained for 90-eV electrons, however Eq. (4) indicates that the equivalent result for 90-eV electrons would be $(2.1 \pm 0.5)\times10^{-4}$ Torr, and this figure compared favorably with the value of $p^* = (1.6\pm1.5)\times10^{-3}$ Torr calculated from Eq. (6) for the limit $\lambda^* = d$. At higher pressure, only electrons from the residual Debye sheath can be detected by the probes for the conditions used. The electrons in the plasma interior are also expelled; however, electrostatic electron-ion coupling results in net plasma motion with high ion energies [Eq. (5)].

It may be concluded that the observed behavior of the high-energy electrons and the ions can be fully explained by an acceleration mechanism based on acceleration of the electrons by nonlinear radiation forces. The 90-eV electron energies that were obtained are attributed to the transfer of quivering energy into translational kinetic energy by the nonlinear ponderomotive force. The different behavioral regimes of nonlinear force acceleration of all the electrons in the laser focus region at very low densities and of combined plasma or ion acceleration at higher densities can be distinguished by the Debye length based on the maximum electron energy.

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Transient Growth in a Current-Carrying Plasma

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It is shown that sheared Alfvén waves in a current-carrying plasma, such as that in a tokamak or in a pinch, may grow linearly in time at a rapid rate. These waves may attain a large amplitude before they eventually decay. The implication of this result is discussed.

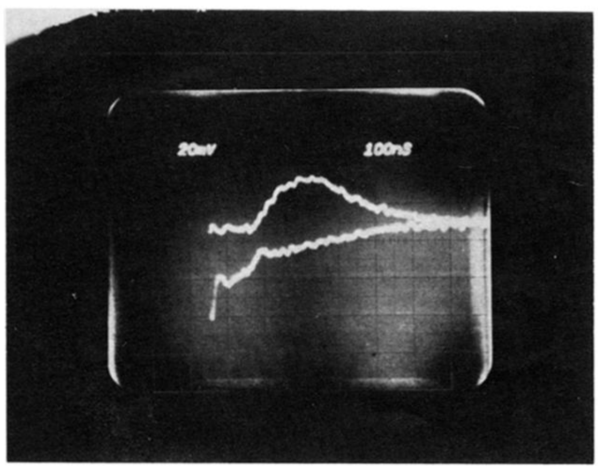
The shear slow of classical hydrodynamics and a current-carrying plasma may be regarded as physical systems with available free energy. In the former case, the energy reservoir is the vorticity, and in the latter case, it is the current. In the presence of dissipation, such free energy may be tapped, and the system would be subject to linear instability. Thus a viscous hydrodynamic shear flow is able to support unstable modes if the Reynolds number is sufficiently high.[1] Similarly, a current-carrying plasma may be subject to the "tearing" instabilities if the plasma medium possesses a small, but finite, resistivity.[2] Indeed, it has been noted [cf. Ref. 2, p. 468] that these two cases are remarkably similar, e.g., Kelvin's cat's eyes appear in both cases.

In the idealized situation where dissipation is absent, there may be no unstable normal modes for either the hydrodynamic shear flow or the current-carrying plasma. Yet both systems may admit perturbations which grow rapidly, reaching a large amplitude before they eventually decay. Such transient growths, in the case of hydrodynamic shear flow, were studied by Orr[3] in 1907 and were regarded by him as the cause of turbulence observed in channel or pipe flows. In this paper, I study the transient growth of perturbations in a current-carrying plasma. The consideration of these transient growths is particularly relevant because, as we shall see, the slab model for this system does not admit unstable eigenmode solutions for virtually any smooth current profile. The conclusion of this work is then that a perturbation may grow linearly in time before it eventually decays. For parameters typical of tokamak geometry, the amplitude may gain by a factor of 50–100 in a time scale of order 5–10 μsec.

As the following analysis would suggest, these transient growths are not restricted merely to a plasma in a tokamak. They may exist as long as there is a longitudinal current in the confined plasma. Their presence may lead to an enhancement of fluctuations, resulting in additional loss of particles or energy. In the worst case, these fluctuations may attain a high enough amplitude that they therefore may cause *other* instabilities (perhaps even major ones) which would not have been excited if there were no such fluctuations.

For simplicity, let us consider a collisionless plasma described by the ideal magnetohydrodynamic (MHD) equations in a slab geometry. Let us assume that the local equilibrium magnetic field is $\vec{B}_0 = B_0(\hat{z} + \hat{x}y/L_s)$, where B_0 is a constant and L_s is a length scale measuring the magnetic shear. If this slab is to simulate the tokamak geometry, the unit vectors $\hat{z}$, $\hat{x}$, and $\hat{y}$ point, respectively, in the toroidal direction, in the poloidal direction, and in the radial direction from the minor axis. The magnetic shear is produced by the longitudinal current, whose density is $\vec{J}_0 = \nabla \times \vec{B}_0 = -\hat{z}B_0/L_s$. This current is assumed to confine a plasma which in equilibrium has a pressure gradient in the $\hat{y}$ direction. The plasma is assumed to possess zero flow velocity in the unperturbed state.

We next consider a small perturbation on such an equilibrium. Let us assume that the plasma



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