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Fast-electron transport in high-intensity short-pulse laser–solid experiments

A R Bell†, J R Davies†, S Guerin† and H Ruhl‡

† Blackett Laboratory, Imperial College, London SW7 2BZ, UK

‡ Theoretische Quantenelektronik, Technische Hochschule, Hochschulstrasse 4A, Darmstadt, Germany

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Abstract. The interaction of short-pulse ($\leqslant$1 ps) high-intensity ($\approx 10^{18}$ W cm^{-2}) lasers with solid targets generates large numbers of energetic ($\approx$100 keV) electrons. The energetic electrons can only penetrate into the solid if the solid can supply an equivalent charge-neutralizing return current. We develop a simple model which shows that in many cases the solid cannot support the required return current and the fast electrons are confined by electric fields to the surface of the target. The target response to laser irradiation depends strongly on the electrical conductivity of the solid.

1. Introduction

Recent developments in laser technology have made it possible to irradiate solid targets with high-intensity ($\sim 10^{18}$ W cm^{-2}) short (<1 ps) laser pulses focused to a laser spot of around 10 μm. It is expected that short-pulse, laser-produced plasmas will develop into an important source of coherent and incoherent x-rays (see the review by Gibbon and Forster [1]). At these high intensities the laser energy is given to high-energy ($\sim$100 keV) electrons which have mean free paths of hundreds of micrometres and a collisional energy loss time of typically a few picoseconds. Consequently, fast electrons can transport the absorbed energy to parts of the target well away from the laser spot over time scales larger than the laser pulselength. The response of the target to laser irradiation is strongly dependent on fast-electron transport which determines the depth and temperature to which the target is heated and the plasma conditions at the surface where the laser energy is absorbed. We show, as discussed previously by Glinsky [2] and others [3, 4], that electric fields can reduce fast-electron penetration into the target to a value much less than the mean free path for energy loss. The penetration depth may depend very strongly on the electrical conductivity of the target. This may explain some curious experimental results.

For the purposes of this paper we assume typically that the laser pulse length is $\tau_{\rm laser} = 1$ ps, the absorbed laser intensity $I_{\rm abs}$ is 10^{18} W cm^{-2}, the laser wavelength is 1 μm, and the laser spot diameter is $2r_{\rm spot} = 30$ μm. A typical target material would be aluminium which has a solid density of $\rho = 2.7$ gm cm^{-3} (electron density $n_{\rm e}$ of 7.8×10^{23} cm^{-3}). The absorbed energy in these conditions is $\epsilon_{\rm laser} = 7.1$ J, and the characteristic fast-electron temperature is $T_0 = 200$ keV [5, 6]. Although the fast electrons are produced by collisionless processes, it is a feature of simulation and experiment that the distribution can, at least in some circumstances, be approximately Maxwellian [7–10].

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With these assumptions, the total number of electrons heated during the laser pulse is 1.5×10^{14}. Their characteristic loss time [11] is $\tau_{ee} = 4.6$ ps (with $\log \Lambda = 8.7$) and the stopping distance ($= \lambda_{ee}/4 = (\tau_{ee}/4)\sqrt{eT_0/m_e}$) is 215 μm. The characteristic time for angular scattering is τ_{ee}/Z ($Z = 13$ for aluminium), and the fast-electron root-mean-square (RMS) range is $R_e \sim (\lambda_{ee}/4)/\sqrt{Z} \sim 60$ μm. The current carried by the fast electrons is $I_{elec} = \epsilon_{laser}/\tau_{laser}/(1.5T_0) = 24$ MA, where T_0 is in electronvolts. If this current were to enter the target as a cylinder of radius r_{spot} of moving electrons, the magnetic field at the surface of the cylinder would be $B_{cyl} = 3200$ MG. If the current penetrated a distance R_e into the target, the energy in the magnetic field would be of the order of $(B_{cyl}^2/2\mu_0)\pi r_{spot}^2 R_e$ $2\log(R_e/r_{spot}) = 5$ kJ. This is energetically impossible implying that such a current cannot be maintained. It must be opposed by an inductively or electrostatically generated electric field which confines the fast electrons near the surface of the target, or else the background thermal plasma must supply a balancing return current. By whatever mechanism it is clear that

$$j_{total} = j_{fast} + j_{thermal} \approx 0 \tag{1}$$

to a good approximation where j_{fast} and $j_{thermal}$ are the currents carried by the fast and thermal electrons, respectively. This implies that not only must the number of thermal electrons entering the laser spot balance the number of fast electrons leaving it, but that everywhere within the target the two currents must be locally nearly in balance.

The continuity equation for fast electrons, combined with equation (1) gives

$$\frac{\partial n}{\partial t} = \nabla \cdot \left(\frac{j_{fast}}{e}\right) = -\nabla \cdot \left(\frac{\sigma}{e} E\right) \tag{2}$$

where $j_{thermal} = \sigma E$ and σ is the conductivity of the thermal plasma. If the fast-electron distribution is Maxwellian and confined by the electric field, $E = -\nabla\phi$, then the fast-electron number density is $n =$ constant $\times \exp(\phi/T_0)$ giving

$$E = -\frac{T_0}{n}\nabla n. \tag{3}$$

Substitution into equation (2) gives

$$\frac{\partial n}{\partial t} = \nabla \cdot \left(\frac{\sigma T_0}{en}\nabla n\right). \tag{4}$$

This is a diffusion equation with a diffusion coefficient $D = \sigma T_0/en$ which is inversely proportional to the fast-electron density.

2. Transport during the laser pulse

Since the electron collisional loss time is larger than the laser pulselength, we can assume that the fast-electron temperature T_0 is constant during the laser pulse. Adiabatic losses can also be neglected because the electric field in the following solution is independent of time. We also assume that the conductivity σ of the thermal plasma is constant and uniform, and that the system is a one-dimensional function of z which is the distance from the surface of the target. With these assumptions the nonlinear equation (4) has a solution

$$n = n_0\left(\frac{t}{\tau_{laser}}\right)\left(\frac{z_0}{z + z_0}\right)^2 \tag{5}$$

where

$$n_0 = \frac{2I_{abs}^2\tau_{laser}}{9eT_0^3\sigma} \qquad \text{and} \qquad z_0 = \frac{3T_0^2\sigma}{I_{abs}} \tag{6}$$

and the fast-electron kinetic energy has been equated to the absorbed laser energy. $I_{\rm abs}$ is the absorbed laser intensity in W m^{-2} and SI units are used throughout except that T_0 is in eV. The total number of fast electrons increases linearly with time which would be consistent with a 'top-hat' laser pulse. The spatial form of the fast-electron distribution (curve a in figure 1) does not vary in time. At the start of the laser pulse the fast-electron density is small so the diffusion coefficient ($\propto 1/n$) is large and the fast electrons easily penetrate the target. During the laser pulse the fast-electron density builds up, reducing the diffusion coefficient and the penetration of those fast electrons that are generated later in the pulse.

The equations for n_0 and z_0 can be recast in the more meaningful forms:

$$n_0 = \left(\frac{I_{\rm abs}}{10^{18}\ \mathrm{W\ cm^{-2}}}\right)^2 \left(\frac{\tau_{\rm laser}}{\mathrm{ps}}\right) \left(\frac{T_0}{200\ \mathrm{keV}}\right)^{-3} \left(\frac{\sigma}{10^6\ \Omega^{-1}\ \mathrm{m^{-1}}}\right)^{-1} 1.7\times 10^{22}\ \mathrm{cm^{-3}} \tag{7}$$

$$z_0 = \left(\frac{T_0}{200\ \mathrm{keV}}\right)^2 \left(\frac{\sigma}{10^6\ \Omega^{-1}\ \mathrm{m^{-1}}}\right) \left(\frac{I_{\rm abs}}{10^{18}\ \mathrm{W\ cm^{-2}}}\right)^{-1} 12\ \mu\mathrm{m}. \tag{8}$$

Milchberg *et al* [12] have measured the conductivity of solid aluminium for temperatures up to 100 eV. For temperatures in the range 10–100 eV the conductivity lies between 0.5×10^6 and $1 \times 10^6\ \Omega^{-1}\ \mathrm{m}^{-1}$. Hence for an aluminium target fast electrons can be expected to penetrate a distance of around 10 μm into the target during the laser pulse. This is much smaller than the collisional range of the fast electrons. It implies that for these conditions the fast electron penetration is limited by the inability of the background plasma to provide a balancing return current. The number of fast electrons generated is so large that they set up a space charge that confines them within around 10 μm of the target surface. Also note that the number density of fast electrons can exceed the critical density and may be a substantial fraction of the solid density. This can be expected to have important consequences for the absorption process since the absorbing plasma may be dominated by fast electrons.

If we further assume that the hot-electron temperature is that given by Beg *et al* [6], $T_0 \approx 100\ [I/(10^{17}\ \mathrm{W\ cm^{-2}})]^{1/3}$ keV, that $I_{\rm abs} = 0.4I$ (40% absorption) and the electrical conductivity is the Spitzer conductivity [11], then equation (8) takes the form

$$z_0 = \left(\frac{T_{\rm cold}}{100\ \mathrm{eV}}\right)^{3/2} \left(\frac{Z}{13}\right)^{-1} \left(\frac{I}{10^{18}\ \mathrm{W\ cm^{-2}}}\right)^{-1/3} 5.9\ \mu\mathrm{m} \tag{9}$$

where $T_{\rm cold}$ is the temperature of the background plasma. The Spitzer conductivity cannot always be expected to apply at these relatively low temperatures. The conductivity of aluminium exceeds the Spitzer value at low temperatures, but the conductivity of other materials can be less and fast-electron penetration severely inhibited.

The ratio of the penetration distance z_0 to the fast-electron collisional range $R_{\rm e}$ (defined above) is smallest at high laser intensity. z_0 decreases with intensity, whereas $R_{\rm e}$ increases since it is proportional to the square of the fast-electron temperature. The laser intensity at which inhibition becomes noticeable can be estimated by equating z_0 and $R_{\rm e}$, in which case we estimate that the electric field restricts electron penetration if

$$I > \left(\frac{T_{\rm cold}}{100\ \mathrm{eV}}\right)^{3/2} \left(\frac{Z}{13}\right)^{-1/2} \left(\frac{\rho}{2.7\ \mathrm{g\ cm^{-3}}}\right) 10^{17}\ \mathrm{W\ cm^{-2}}. \tag{10}$$

It appears that inhibition might occur at low laser intensities if $T_{\rm cold}$ is low, but at these lower temperatures the Spitzer conductivity and hence this estimate of the critical intensity is invalid. Nevertheless, it implies that the effect might be measurable in a wide range of experiments provided that the chosen target has a low electrical conductivity.

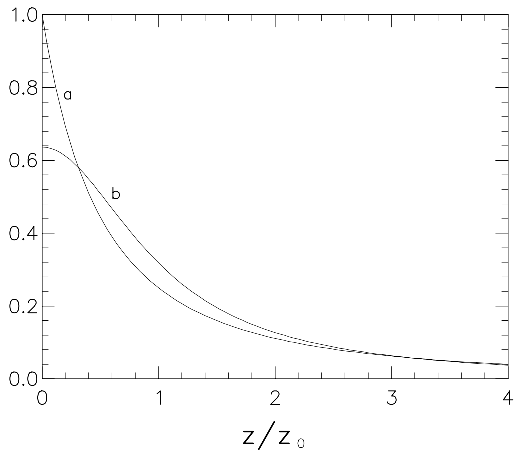

Figure 1. The two solutions for the spatial dependence of the fast-electron number density: curve a, during the laser pulse: $[z_0/(z+z_0)]^2$ (equation (5)); curve b, immediately following the laser pulse: $(2/\pi)z_0^2/(z^2+z_0^2)$ (equation (12)).

Equation (8) suggests that fast-electron penetration depends strongly on the target material. Very different fast-electron penetration might occur for different materials. The ability or inability of fast electrons to escape the laser spot will affect the plasma conditions there and other processes such as absorption can be expected to vary between targets of different materials.

3. Transport after the laser pulse

Equation (5) gives an appropriate solution to equation (4) while the fast electron population is being produced by laser irradiation. The electrons are confined to the surface of the target during the laser pulse, but they may diffuse deeper into the target at later times. Adiabatic losses would then be important because the fast electrons would be contained within an expanding potential well. We seek a self-similar solution of equation (4) in which the spatial shape keeps the same form but expands in time with a scalelength $L(t)$. Assuming that there is no further heating of fast electrons, and neglecting collisional losses again, energy conservation implies an adiabatic decrease in temperature

$$T(t) = \frac{T_0 z_0^{2/3}}{L(t)^{2/3}} \tag{11}$$

where z_0 is the scalelength when $T = T_0$, which is the temperature at $t = \tau_{\text{laser}}$ when the laser pulse switches off. We assume that it is still acceptable to treat the electrons as Maxwellian with a temperature T which is uniform spatially but varies with time.

Equation (4) has a self-similar solution which conserves total electron number:

$$n = \frac{2n_0 z_0}{\pi}\frac{L}{z^2+L^2} \tag{12}$$

$$L(t) = z_0\left[\frac{5\pi\sigma T_0}{3en_0 z_0^2}(t-\tau_{\text{laser}})+1\right]^{3/5}. \tag{13}$$

The constants of integration have been chosen so that $L(\tau_{\text{laser}}) = z_0$ and the total number of electrons is equal to that given by the density profile in equation (5) at time $t = \tau_{\text{laser}}$.

As shown in figure 1, the form of the spatial dependence is not the same in equations (5) and (12), but the forms are sufficiently close for us to build up a semiquantitative picture of how the fast-electron distribution evolves both during and after the laser pulse but before collisional energy losses become important.

Using the expressions given for n_0 and z_0 in equation (6), equation (13) simplifies to

$$L(t) = 1.78 z_0 \left[\frac{t}{\tau_{\mathrm{laser}}} - 0.618 \right]^{3/5}. \tag{14}$$

For times greater than a few times the laser pulselength, the penetration distance increases as $L \propto t^{3/5}$. From equation (11) the fast-electron temperature decreases as $T \propto t^{-2/5}$. For the typical parameters chosen here $\tau_{\mathrm{laser}} = 1$ ps, and the fast electron loss time is $\tau_{\mathrm{ee}} = 5$ ps giving a stopping time of $\tau_{\mathrm{ee}}/3 \approx 1.7$ ps for a fast electron with energy eT_0. Hence this phase of adiabatic expansion will not last long and the fast electrons do not penetrate a distance much greater than z_0. In other cases adiabatic expansion might allow the fast electrons to penetrate much further into the target, especially if the irradiance is high or the target has a low Z.

4. Discussion

Experiments at high intensity are often simulated with collisionless particle-in-cell or Vlasov codes in which the starting plasma temperature is in the kiloelectronvolt range [13–15]. A collisionless high-temperature plasma has a high conductivity. Hence the return current is much more easily provided than in real experiments in which the thermal target material is at a relatively low temperature and highly collisional. Consequently there must be some doubt over the realism of many simulations of high-intensity laser–solid experiments.

Inhibition of fast-electron transport due to space-charge-induced electric fields was demonstrated by Bond *et al* [4, 16] more than a decade ago at the lower laser intensity of 3×10^{15} W cm^{-2} using low-density foam targets. Experiments at intensities of 10^{18} W cm^{-2} have only become possible recently. As yet there have been few systematic investigations of fast-electron transport. Luther-Davies *et al* [17] performed experiments with laser pulselengths which were relatively long (20–100 ps) by recent standards. Their highest intensities were 3×10^{17} W cm^{-2}. They measured the Kα emission from nickel layers buried at varying depths. Their space-resolved images of the emission showed that fast electrons penetrated further into the target in the region outside the laser spot than in the part of the target underlying the laser spot. This may be due to magnetic insulation [18–22] as they suggest, but it could also be due to electric fields having a stronger effect where the fast-electron flux is greatest.

Jiang *et al* [23] have experimented with layered targets when the laser intensity was a few times 10^{18} W cm^{-2} at a wavelength of 0.5 μm. The pulselength was 300 fs. They found that the fast electrons penetrated through approximately 3 μm of aluminium. They calculate that this small penetration distance would be consistent with a fast-electron temperature of only 25 keV, but they do not report confirmation of this low temperature by an independent measurement. Other experiments at this intensity have measured fast-electron temperatures in excess of 100 keV [5] which would have a collisional mean free path more than 16 times that of electrons at 25 keV (mean free path $\propto T^2$). Furthermore, 25 keV is an order of magnitude smaller then the electron oscillation energy in the laser field which usually characterizes the fast-electron temperature [14]. It is therefore conceivable that the small electron penetration measured by Jiang *et al* is due to inhibition by electric fields.

A similar effect is found by Teubner *et al* [24] for experiments with a laser wavelength of 0.25 μm and a peak intensity of 5×10^{18} W cm^{-2}. From the measurement of the penetration distance they deduce a fast-electron temperature of 8 keV which is much less than expected. This may be further evidence for transport inhibition by electric fields.

Beg *et al* [6] deduce a hot-electron temperature of between 70 and 200 keV from buried $K\alpha$-emitting layers, but find evidence from the x-ray bremsstrahlung emission that the temperature of the hot component is 390 keV. They advise caution in the interpretation of both these temperature measurements, but note that 'resistive inhibition' might be reducing the electron range giving an incorrectly low $K\alpha$ estimate of the hot-electron temperature. They also find some suggestion that the fast-electron penetration is material dependent.

The absorption of short intense laser pulses has been more thoroughly investigated than transport. The results differ widely between experiments. At intensities around 10^{18} W cm^{-2} Beg *et al* [6] measure absorption in the region of 40% and Schnurer *et al* [25] up to 50% absorption, yet Price *et al* [26] measure 10%. The difference in these various results may be due to differences in laser spot size or spot quality, target preparation or laser prepulse, but the efficiency of fast-electron transport out of the laser spot may also have an effect on absorption. If the fast electrons are contained near the target surface the laser radiation could be interacting with a much hotter plasma than if the fast electrons are free to escape.

Over the next few years we can expect high-intensity short-pulse laser–solid experiments to become more reproducible as the technology and theoretical understanding increases. The aim of this paper is to show that the ability of the background thermal plasma to supply a return current may be an important factor in determining experimental behaviour. Penetration into the solid target may be much reduced below the expected collisional range of fast electrons.

Space-charge inhibition of fast-electron transport may also have implications for the fast ignitor approach to inertial confinement fusion [27]. In this approach a high-intensity, short-pulse laser delivers energy in the form of fast electrons to the compressed core. The fast electrons penetrate the core, heating it, and initiating fusion reactions. Electrical inhibition of fast-electron transport may require the laser energy to be absorbed deeper into the dense core.

Acknowledgments

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