Hot electron generation at a steep interface in super intense laser-matter interaction

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Abstract. Super intense laser light (> 10²⁰ W/cm²) is able to sweep preplasma over short times and compress the preplasma density gradient typically generated by the prepulse of today’s high-intensity, high energy laser systems. Hot electron generation at steep plasma density gradients has been studied in a previous paper [Kemp et al., Phys. Rev. Lett. (2008)], which identified a mode of hot electron acceleration that is characterized by the formation of low-density shelf in front of the target. In this paper, we study the direct interaction mode, which appears after the shelf plasma interaction. We propose an analytical model to describe the absorption, explain electron energy spectra, and identify the parameter regime where the direct interaction mode becomes important.

1. Introduction
Current advanced laser technology is capable of producing high contrast ratio (10¹⁰) laser pulses with relativistic intensity (> 10¹⁸ W/cm²) and ultra short duration (few tens of fs). Such high contrast short laser pulses form only a small pre-plasma thus can directly interact with steep interface of targets, which is important because of various potential applications e.g. isochoric heating of dense plasmas [1], inertial confinement fusion, bright hard and soft x-ray sources [2]. Although strong bulk heating beyond 100 eV is expected in this direct heating mode, the underlying heating mechanism is not fully understood.

In an intense laser solid interaction, there are basically three different stages in a one-dimensional interaction scenario. The first important acceleration is the ponderamotive acceleration, near critical density, n_c, according to Wilks’s scaling [3]. Then a long and intense laser pulse can sweep off the preformed plasma, and form a shelf plasma with the density ~ n_c in front of the steepened interface [4]. Finally the shelf plasma density is being reduced through the interaction, and drops much lower than n_c. Following that, the laser light interacts with dense plasma directly at the steepened interface. In this paper we study this direct interaction mode.

2. Numerical study of LPI at steep density gradient
In this section we demonstrate a one-dimensional simulations, for laser-matter interaction which focus on direct interaction mode, using our PICLS code [5]. Hot electrons produced in this mode are important for energy transport inside a solid target. A slab target with fully ionized deuterium plasma extends from 20 - 40 λ in a 40 λ simulation box. We used a computational
resolution of $\Delta x = 1\mu m/78$ and $\Delta t = \tau/78$ to resolve the highest density, $\sim 400n_c$. Each mesh ($\Delta x$) contains 300 deuterium ions and 300 electrons. A linearly polarized laser pulse ($\lambda = 1\mu m$) with semi infinite envelope and ($a = 10$), rising up in 15 $\tau$ with a gaussian profile, incidents normally on slab targets of width $20\lambda$. Absorbing boundaries are used for particles and electromagnetic fields. To account plasma density dependence on hot electron generation and their energy scaling we performed a set of simulations with different slab targets of densities ranging between 40 to 400 $n_c$. Simulations were carried out for approximately 500 fs (150 $\tau$). Figure 1 shows time evolved electron phase plots at the irradiation surface during one laser oscillation period. Electrons are trapped in electrostatic potential, at target surface, and form a blob structure of thickness $\sim 3l_s = 3c/\omega_p$. Peak longitudinal momentum of this blob depends on electrostatic potential, which is balanced by laser photon pressure [4]. Note here that the total number of electrons trapped in this blob are almost constant during the interaction. Leaving electron jets are balanced by return current which supplies electrons to the blob. In next sections we will introduce analytical models for the electron acceleration in the blob. In Fig. 1, (b), (c) and (d) we see the continued oscillation of trapped electrons and $2\omega$ electrons extraction from outer blob region by $J \times B$ mechanism [6] in one laser period and then injected back into the target. Except $2\omega$ jets, additional 5-6 electron jets with lower energy are observed emerging from the rear end of blob. We call them DC-ponderomotive (DCP) electrons, as they are accelerated via DC field of laser as discussed later. The energy of the DCP electrons depends on the peak momentum of the blob electrons. Note here that the electrons in the central region are always trapped and never come out from the blob, so the electrons forming $2\omega$ and DCP electrons are originally located deeper in the target and they are introduced to the outer most orbit of the oscillating blob by means of the return current.

![Figure 1](image1.png)

**Figure 1.** Electron phase evolution, $x - p_x$. (a) Electrons emerging, arrows show electron flow (b) 1st jet of $2\omega$ extraction (c) $2\omega$ and DCP electrons at later time (d) 2nd jet of $2\omega$. Arrows point in the direction of rotating blob in phase space. Dashed line shows the position of peak electron density.

For $200n_c$ target, we observed a similar blob structure of thickness $\sim 3l_s$ and DCP electrons jets, but we do not see $2\omega$ jet. However the height of blob is reduced by a factor of $\sqrt{1/n_c} = \sqrt{40/200} = 0.44$. Again, the maximum number of DCP electrons group is five to six in one laser period, similar to $40n_c$ density case. The number of DCP electron jets has a weak density dependence [7]. In other density cases such as, $70n_c$, $100n_c$, $300n_c$ and $400n_c$, except $70n_c$ which has an insignificant $2\omega$ jets, we do not observe $2\omega$ electrons. Electron energy
spectra obtained for different target densities are plotted together in Fig. 2. This picture gives a clear evidence for disappearance of high energy tail in spectrum due to the $2\omega$ jet suppression.

3. Model for direct interaction: DC-ponderomotive (DCP) electrons

A normally incident laser on target exerts a pressure $P_L$, $P_L = (1 + \eta)I/c$ where $\eta$ is reflectivity and $I$ is laser intensity. This pressure, which results due to the non oscillating component of ponderomotive force, pushes electrons more than the ions, hence a charge separation is established and an electrostatic potential $\phi$ is excited at the target surface [7]. This potential gap can be estimated by conserving the momentum flux of mass flow (plasma pressure) with the light pressure, $(1 + \eta)I/c = n_e e\phi$, at the reflection surface [4],[8]. If electrons accelerate in this potential well up to energy $e\phi$ and gain longitudinal momentum $P_e$, then their energy is,

$$\epsilon = m_e c^2 \left( \sqrt{1 + P_e^2 / m_e^2 c^2} - 1 \right) = (1 + \eta)I/n_c c,$$

(1)

here $P_y$ is neglected, $m_e$, $c$, $n_e$ are electron mass, light velocity and electron density, respectively. In simulations we observed that the DCP electrons escaping from the oscillation region have peak energy which is the average energy of the oscillations, then the energy of DCP electrons is given by $\langle \epsilon \rangle = m_e c^2 a^2 (1 + \eta)n_e / 4n_e [1]$, here, time averaging in one laser period gives $< P_e^2 > \approx < P_e^2 > / 2$ and using relations $I/c = m_e c^2 a^2 n_e / 2$ and $a = eE_L / m_c c\omega$. The peak energy of oscillations in the blob for $40n_c$, $70n_c$ and $300n_c$ density cases are calculated by Eq. (1) and indicated with arrows in Fig. 2, which are consistent with the observed spectrum. For example, in $40 n_c$ case the arrow clearly shows the transition point between DCP electrons and $J \times B$ electrons energy in the spectrum.

Fig. 1 shows that the electrons are trapped at the target surface and form a blob like structure of thickness $\sim 3l_s$. Note here that area of blob in phase space is conserved, meaning no net accumulation of charge or energy, particle conservation gives a condition; $v_d n_d \simeq v_r n_e$, here $v_d, n_d$ and $v_r, n_e$ are the speed and density of DCP and return current electrons, respectively. Energy flux conservation gives us the condition $I_{in} - I_{rf} = \chi I_{in} \simeq \epsilon_d v_d n_d$, here $\chi$ is absorption coefficient and $\epsilon_d$ is energy of the DCP electrons. The energy of electrons in return current, which scales with $(n_d / n_e)^2$, is negligible compared to that of the DCP electrons. Plugging in the DCP electron velocity $v_d$ in energy flux conservation condition calculated from the expression which gives average energy of DCP electrons.

$$\chi \simeq (2/a) (n_d / n_e) \cdot (1 - 1/\gamma_d) \sqrt{n_e / n_e},$$

(2)

here $\gamma_d = \sqrt{1 + < P_e^2 > / m_e c^2} \sim 1 + a^2 (n_e / n_e)$, and $(1 + \eta)/2 \sim 1$. The factor $(1 - 1/\gamma_d)$ in Eq.(2) is very small for $a^2 (n_e / n_e) \ll 1$. To estimate maximum absorption we solve $\partial \chi / \partial a = 0$ and find $\chi_{max} \simeq 0.6 (n_d / n_e)$ at $a \simeq 1.27 \sqrt{n_e / n_e}$, which indicates that the absorption by DCP electrons becomes efficient for $a > 10$ and $n_e > 100n_c$. A rough estimation of DCP electron density, is $n_d \sim n_e / 20$ for $400n_c$ and $\sim n_e / 10$ for $100n_c$. Though this is a small fraction of input energy but DCP electrons deposit energy effectively inside the target in a small volume, which is important to create keV energy - solid density plasmas.

In Fig. 3 we plot Eq. (1) (DCP electrons), and $J \times B (2\omega)$ scaling obtained from Helmholtz boundary solution [9]. Electron momentum observed in simulations agrees well with $2\omega$ scaling below $100n_c$, while with the DCP energy scaling above $100n_c$. A transition occurs around $100n_c$, in dominating mechanism to produce the maximum energy electrons. $2\omega$ jets are disappeared in higher density targets ($>100n_c$) since they do not have sufficient energy to overcome potential gap so remain trapped, also their suppression reduces the total absorption. Since a circularly polarized laser pulse does not produce $2\omega$ electrons, then the absorption will occur mainly because of DCP electron acceleration. To see laser polarization effect on DCP electron we
performed simulations with circularly polarized laser light while keeping other laser and target parameters similar. Note here that in simulations we have found that the ion contribution to total absorption is extremely small, such as < 1% for \( a=10 \), and it drops to \( \sim 0.001\% \) for \( a=1 \). Hence the absorption is mainly by the electrons. The ion absorption rates are consistent with the Denavit [10]. We measured the total absorption for the period of 150 \( \tau \) (\( \sim 500\) fs). In Fig. 4, we can see that the absorption is low \( \sim 1\% \) and almost constant with \( a = 1 \), since the DCP electrons are inefficient. Also, the absorption for circularly polarized laser light is identical to the absorption of linearly polarized light, above \( 100n_c \). We can conclude here that for density \( > 100n_c \), DCP remains the only absorption mechanism.

4. Summary

We studied electron acceleration at overcritical plasma surface and demonstrated the production of DC-ponderomotive (DCP) electrons by non-oscillating laser photon pressure. We have also proposed an analytical model of ‘ponderomotive’ electron acceleration and demonstrated that the injection frequency and energy of DCP electrons are independent of laser polarization. More importantly, the DCP mechanism for electron acceleration dominates in high density targets, such as \( n_c \geq 100n_c \). Absorption by DCP electrons does not depend on laser polarization above \( n_c > 100n_c \). These DCP electrons, with short stopping ranges and relatively high density, \( \sim 10\% \) of the target density, have a great potential to heat the target in a limited volume. We derived optimal conditions for the absorption via DCP electrons, which suggest, i.e., in targets of density 100 \( n_c \) or above the contribution from DCP electrons to the total laser absorption will be maximum for a pulse with \( a \sim 10 \).

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