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Charge separation effects at the solid-vacuum interface and ion acceleration with two electron populations

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The electrostatic field at the solid-vacuum interface produced by two equilibrium electron populations with different temperatures is analytically calculated and investigated in terms of plasma parameters. The hot and the cold electron populations represent the laser accelerated electrons (typical energy of a few MeV), and the ambient electrons carrying the return current and heated due to the finite resistivity of the solid target. In the frame of this model, the ion acceleration process in laser-solid experiments is discussed. It is compared with the numerical results obtained by using a Boltzmann-Vlasov-Poisson model of ion acceleration.

Among the various physical processes arising during the interaction between a relativistically intense laser pulse and a solid thin film, the acceleration of light and heavy ions at the rear side of the target is of particular interest, both from a fundamental point of view, and for the many potential applications. A clear understanding of the relevant phenomena involved is therefore mandatory in order to fully exploit this mechanism of ion acceleration. In particular, it is important to describe the properties of the extremely intense accelerating electric field produced, at the rear side of the target, as a result of the expansion of the relativistic electron cloud created at the front surface by the laser pulse.

We have investigated analytically the electrostatic field, in a one-dimensional geometry, by assuming that two different equilibrium electron populations are present at the rear surface: a "hot" population, representing the laser accelerated electrons (with a typical energy of the order of a few MeV), and a "cold" population, representing the ambient electrons carrying the return current and heated due to the finite resistivity of the solid target [1]. The cold ions are localized in the $x<0$ region, $x=0$ being the solid-vacuum interface and where the x axis is directed normally to the target surface. The electrostatic potential ϕ satisfies the Poisson-Boltzmann equation:

$$\partial_x^2\phi = 4\pi e\left[n_{0h}e^{e\phi/T_h} + n_{0c}e^{e\phi/Tc} - \theta(-x)Zn_{0i}\right] \qquad (1)$$

which, by imposing neutrality for $x \to -\infty$ and vanishing electric field and particle number for $x \to +\infty$, admits the implicit solutions [1]

$$\int_{\varphi(0)}^{\varphi(x)} \frac{d\varphi}{\left[\exp\varphi + ab\exp(\varphi/b) - (1+ab) - (1+a)\varphi\right]^{1/2}} = -\sqrt{2}\,\frac{x}{\lambda_{dh}} \quad (2)$$

$$\int_{\varphi(0)}^{\varphi(x)} \frac{d\varphi}{\left[\exp\varphi + ab\exp(\varphi/b)\right]^{1/2}} = -\sqrt{2}\,\frac{x}{\lambda_{dh}} \quad (3)$$

inside and outside the target, respectively. In the above equations, $n_{0h(0c)}$ and $T_{h(c)}$ are the unperturbed hot (cold) electron density and the hot (cold) electron temperature, respectively, n_{0i} and Z the ion density and charge state, $\theta(x)$ the Heaviside function. Moreover, $\varphi \equiv e\phi/T_h$ is the dimensionless potential, $\varphi(0)=(1+ab)/(1+a)$, with $a\equiv n_{0c}/n_{0h}$ and $b\equiv T_c/T_h$, and $\lambda_{dh}=(T_h/4\pi n_{0h}e^2)^{1/2}$ is the hot electron Debye length. The consequences of the adopted boundary conditions are discussed elsewhere [2]. The maximum electric field, in $x=0$, reads

$$E(0) = \sqrt{2}\,\frac{T_h}{\lambda_{dh}}\left\{\exp\left[-b\left(\frac{1+ab}{b+ab}\right)\right] + ab\exp\left(\frac{1+ab}{b+ab}\right)\right\}^{1/2} \quad (4)$$

The cold-to-hot electron pressure ratio $ab=p_{0c}/p_{0h}$ strongly influences this value. In typical experimental conditions, $a>>1$ and $b<<1$; then, if the hot pressure is greater than the cold one, that is, $ab<<1$, $E(0)$ is approximately given by $E(0) \approx \sqrt{2}\, T_h/e\lambda_{dh}$, while in the opposite case, $ab>>1$, we can approximate $E(0)$ as

$$E(0) \approx \sqrt{\frac{2}{\mathrm{e}}ab}\,\frac{T_h}{e\lambda_{dh}} \quad (5)$$

where e is the Neper's number. Cold electrons determine the distribution of the electric field inside the target; indeed, from Eq. (2), it turns out that it can penetrate over a distance of the order of a few cold Debye lengths $\lambda_{dc}=(T_c/4\pi n_{0c}e^2)^{1/2}$, as shown in Fig.1. Notice that the penetration of the electric field directly influences the number of ion-layers accelerated and the corresponding energetic spectrum.

Our investigation demonstrates the important role played by the cold electron parameters in the determination of the properties of the accelerating electric field (maximum value and spatial distribution). While the hot electron parameters depend essentially on the interaction between the laser pulse and the target at the front surface, those of the cold electrons are determined by the target properties. In order to get a reasonable estimation of T_c, we have developed an analytical model for the resistive Ohmic heating of the cold electron population, solving the Fourier equation for T_e [3], neglecting the target thermal conductivity and assuming generic power-law dependences for the target electrical resistivity $\eta = \eta_k(T_e/T_k)^\alpha$ and for the cold electron heat capacity $C_e(T_e) = C_k(T_e/T_k)^\beta$, where η_k and C_k are the electrical resistivity and the electron heat capacity at T_e=T_k, respectively.

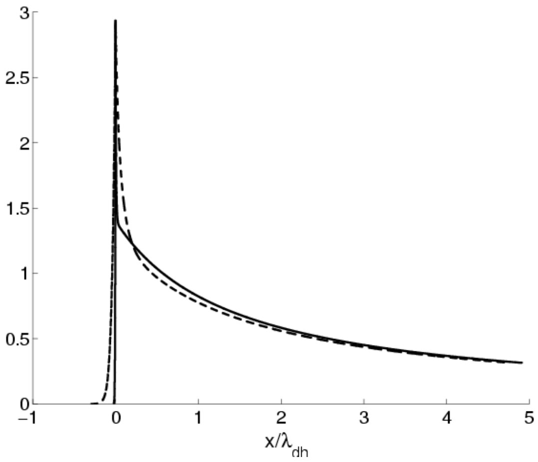

Fig. 1 – Spatial distribution of the normalized electric field $eE\lambda_{dh}/T_h$ for *ab*=10 and *b*=0.01 (solid line) and 0.1 (dashed line)

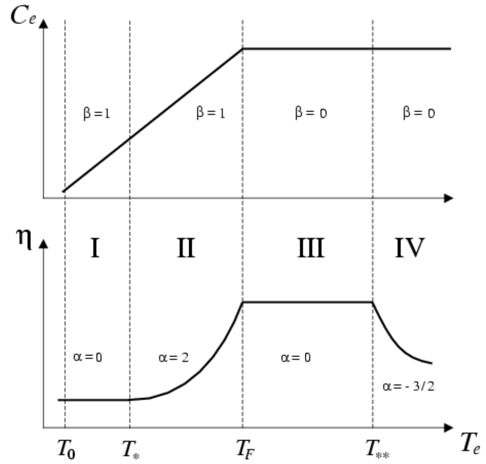

Fig. 2 – Temperature dependence of the electron heat capacity and resistivity: T_0 is the room temperature and T_F is the Fermi temperature

By choosing T_k as the temperature at the generic initial time t_k, the electron temperature T_e is given by [1]

$$T_e(t) = T_k\left[1+\frac{\eta_k j^2}{\gamma C_k T_k}(t-t_k)\right]^\gamma \tag{6}$$

for α -β <1, where $\gamma=(1-\alpha+\beta)^{-1}$ and j is the current density, while if α -β =1 we have

$$T_e = T_k \exp\left[\frac{\eta_k j^2}{C_k T_k}(t-t_k)\right]. \tag{7}$$

Different values of the coefficients α and β have been used, depending on the instantaneous value of T_c, as shown in Fig. 2 (see Ref.1 for details), in order to describe the different physical regimes produced during the heating process, which holds on for a time of the order of the pulse duration. In Fig. 2 T_* corresponds to the temperature where electron-electron collisions become important and T_{**} is the minimum temperature for the Spitzer conductivity to be valid. The model has been applied to the case of a target of Al, for which final values of T_c of the order of 100 eV for short pulses (tens fs) and 1 keV for longer pulses (hundreds fs or more) are estimated, respectively.

The electric field calculated on the basis of Eqs. (2) and (3) has been considered as the initial condition in a numerical hybrid Vlasov-Boltzmann-Poisson model (collisionless kinetic ions and isothermal Boltzmann electrons coupled by the electrostatic field), used to study the acceleration of light and heavy ions at the rear side of the target. Spatial profiles, energy distributions, and maximum energies of accelerated ions have been analyzed as functions of the system parameters [4]. Fig. 3 shows the time-averaged proton energy spectrum $dN_p/d\varepsilon$, accelerated from a CH layer, for different values of the pressure ratio *ab*: 0.1; 1; 2; 10. The numerical solutions are stopped at the time t=10t_0, t_0 being equal to the characteristic proton acceleration time $\lambda_{dh}/(T_h/m_p)^{1/2}$, where m_p is the proton mass. In these simulations the free parameter is the cold electron temperature. It results that by increasing T_c, the spectrum becomes broader and the maximum final ion energy larger. This is in qualitative agreement with our analytical model: as T_c increases, λ_{dc} also increases and the field penetrates more deeply inside the target (see Fig. 1), accelerating a larger number of protons, localized in a deeper spatial region, which eventually results in a broader accelerated energetic spectrum.

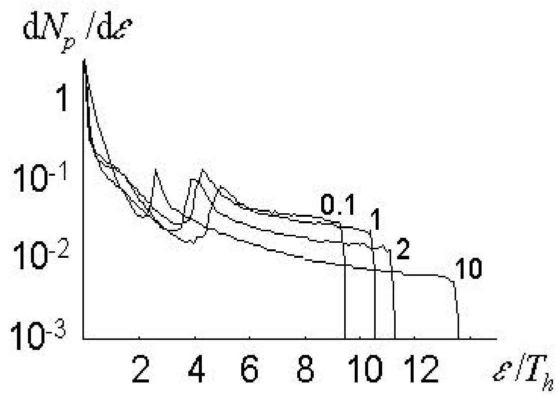

Fig. 3 – Dependence of the spectrum of accelerated protons on the cold-to-hot electron pressure ratio at time 10 t_0

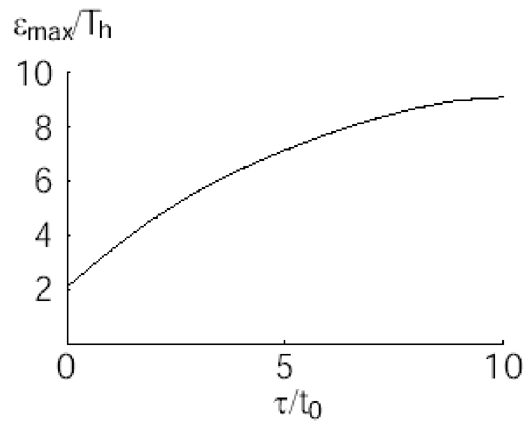

Fig. 4 – Dependence of the maximum energy of fast protons on the laser pulse duration (modeled by the cooling time of hot electrons)

Moreover, if we consider the spectrum at the same final time, the maximum ion energy increases with T_c because the initial electric field is larger, according to Eq. (5). In these simulations the number of protons accelerated is lower than the number of hot electrons in the cloud, so in a first approximation the properties of the electric field are not dramatically altered by the ion dynamics.

The Vlasov-Boltzmann-Poisson model allows to analyze the efficiency if the ion acceleration in function of the laser pulse duration. Figure 4 shows that the pulse duration of the order of 10 t_0 is needed in order to acceleration ions to laser energies.

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