# Beam Instabilities in Laser-Plasma Interaction: Relevance to Preferential Ion Heating

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We propose a new mechanism for anomalous ion heating in ultraintense laser plasmas. This mechanism is based on the excitation of an electron two-stream instability that is driven by the fast electron beam that resonantly decays into ion-acoustic waves. These low frequency waves are then strongly damped by the ion collisions in the dense plasma. The model gives a simple explanation for the preferential heating of the bulk ion population observed in recent laser experiments in the petawatt regime. In particular, this work provides an explanation for the different energy loss in the Au and CD plasmas, in cone-guided fast ignition experiments.

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There is now a growing body of evidence of anomalous ion bulk heating in laser-plasma interaction experiments, when the laser power approaches the petawatt (PW) regime. In integrated fast ignition experiments at Osaka University, the PW laser energy transfer to the compressed plasma was measured to be $\sim$20 percent [1,2]. This is much higher than expected, given the areal density of the compressed plasma of 100–250 mg cm$^{-2}$ and the classical range of MeV electrons in the generated fast particle beam. At the Rutherford Appleton Laboratory, the number of MeV electrons reaching an electron spectrometer located on the chamber wall was reduced by over an order of magnitude when they had to traverse a compressed carbon-deuterium (CD) plasma with an areal density of 20 mg cm$^{-2}$ in cone-guided experiments [3]. It appears from these measurements that an anomalous stopping and/or scattering mechanism is present in the interaction.

Recently, Campbell *et al.* [4] simulated the ion heating in the Osaka University experiments and noted that the ion heating was maximized close to the boundary between the compressed core and the tip of the cone. This was attributed to fluctuating magnetic fields creating an electric field that caused the fast electrons to execute chaotic orbits, thereby increasing the plasma resistivity.

Here we propose an alternative explanation for this anomalous resistivity. It is based on the existence of two coupled processes. First, the fast electrons created by the laser pulse interact with the resulting return current [5] and produce an intense electrostatic field, by means of a two-stream instability. Second, the resulting electrostatic waves become modulationally unstable and decay resonantly into ion-acoustic waves. This mechanism is not efficient for electron heating but leads to an efficient ion heating which could explain the observations. Anomalous stopping of the fast electron beam by electron plasma turbulence was recently discussed [6], but the coupling with the ion oscillations was not considered.

The energy cascading in the solid target takes the following form: the intense laser pulse hits the surface of the solid target and transfers a significant amount of its energy to a relativistic electron beam. This beam of energetic electrons produces a return current and the resulting counter-streaming electron beams excite an electron two-stream instability. The resulting plasma waves have relativistic phase velocities and can easily become modulationally unstable by decaying into heavily damped ion-acoustic oscillations. This results in the occurrence of anomalous ion heating process which efficiently dissipates the energy transferred to the ion oscillations.

Alternatively, ion heating could result from the occurrence of Weibel instabilities and the consequent filamentation of the fast electron current [7], but there is strong indication from particle-in-cell simulations that filamentation is not important for thick targets and only occurs very close to the target surface. The conversion of fast electron beam energy into magnetic field energy is therefore quite small [8]. Another possible explanation could be the excitation of ion oscillations directly by the return current but, as shown later, the electron velocity of the return current is usually very large and, for this reason, it is not well suited to excite low frequency ion oscillations that would transfer their energy directly to the ions. These two more direct heating processes seem to be less promising to explain the observations.

In order to establish our two-step model for anomalous ion heating, we consider the various beam instabilities that can occur in laser produced plasmas and establish their relevance to ion heating in the fast ignition scenario. We will restrict our discussion to the electrostatic case. This will involve ion and electron instabilities in a plasma with the following constituents: (i) a fast electron beam, (ii) a return current, (iii) an ion background and (iv) a broadband plasmon spectrum. In our analysis, we will extend our previous work [9] in several directions, namely, by including relativistic particle beams, instability growth of the plasmon turbulence and particle collision frequencies.

Initially, our plasma is made of a background population of ions, a fast electron beam and a return current. Assuming that the net current is nearly equal to zero, we can establish a relation between the velocities of the two electron populations $\vec{J} = -e(n_{0e}\vec{u}_e + n_{0f}\vec{u}_f) \simeq 0$, where the subscripts $e$ and $f$ refer to the return current and to the fast electron beam, respectively. We can conclude that $u_e \simeq (n_{0f}/n_{0e})u_f$, with $u_f \simeq c$.

Because of the electron two-stream instability, a broadband plasma turbulence is excited. The plasmon gas can be described by a kinetic equation, which is the equation of conservation of the number of plasmons, plus a source term, as given by

$$\left(\frac{\partial}{\partial t} + \vec{v}_{k'} \cdot \frac{\partial}{\partial \vec{r}} + \vec{F}_{k'} \cdot \frac{\partial}{\partial \vec{k}'}\right) N_{k'} = 2\Gamma_{k'} N_{k'} \qquad (1)$$

where $N_{k'} = W_{k'}/\hbar\omega_{k'}$ is the plasmon occupation number, where $W_{k'}$ is the electrostatic energy density, $\omega_{k'} = (\omega_{pe}^2 + S_e^2 k'^2)^{1/2} \simeq \omega_{pe}$ is the plasmon frequency, and $S_e = (3T_e/m_e)^{1/2}$ is the electron thermal velocity. We have also used the plasmon growth rate $\Gamma_{k'}$, the group velocity $\vec{v}_{k'} = S_e^2 \vec{k}'/\omega_{k'}$ and the force $F_{k'} = -(e^2/2\epsilon_0 m_e \omega_{k'})\nabla n_e$, where $n_e$ is the electron plasma density.

Let us now study the plasma stability in the presence of the plasmon field that results from the electron two-stream instabilities. We have to use the relativistic fluid equations for the fast electrons, the return current and the plasma ions ($\alpha = f, e, i$)

$$\begin{aligned} &\frac{\partial n_\alpha}{\partial t} + \nabla \cdot n_\alpha \vec{v}_\alpha = 0, \\ &\frac{\partial \vec{p}_\alpha}{\partial t} + \frac{\vec{p}_\alpha \cdot \nabla \vec{p}_\alpha}{m_\alpha \gamma_\alpha} = -q_\alpha \nabla \psi - \frac{\nabla P_\alpha}{n_\alpha} - \nu_\alpha \vec{p}_\alpha, \end{aligned} \qquad (2)$$

where $n_\alpha$, $\vec{v}_\alpha$, and $P_\alpha$ represent the density, velocity, and pressure of the various particle populations $\alpha$, and $\nu_\alpha$ their collision frequencies. We also have $\gamma_\alpha = (1 - v_\alpha^2/c^2)^{-1/2}$. In this equation we have also used the electrostatic potential $\psi$, determined by the Poisson's equation $\nabla^2\psi = -(1/\epsilon_0)\Sigma_\alpha q_\alpha n_\alpha$, where $q_\alpha$ are the particle charges, $q_e = q_f = -e$, and $q_i = Ze$.

We now consider density and velocity perturbations $\tilde{n}_\alpha$ and $\delta\vec{v}_\alpha$, such that $n_\alpha = n_{0\alpha} + \tilde{n}_\alpha$ and $\vec{v}_\alpha = \vec{u}_\alpha + \delta\vec{v}_\alpha$. This also implies that $\vec{p}_\alpha = \vec{p}_{0\alpha} + \delta\vec{p}_\alpha$ and $\gamma_\alpha = \gamma_{0\alpha} + \tilde{\gamma}_\alpha$. Replacing in the above equations, and using the simplifying notation $D/Dt = (\partial/\partial t + \vec{u}_\alpha \cdot \nabla)$, we can obtain the equation for the perturbed densities

$$\left(\frac{D}{Dt} + \nu_e\right)\frac{D}{Dt}\tilde{n}_e + \frac{\omega_{p\alpha}^2}{\gamma_{0\alpha}^3}\sum_\beta \frac{q_\beta}{q_\alpha}\tilde{n}_\beta - \frac{S_\alpha^2}{\gamma_{0\alpha}^3}\nabla^2\tilde{n}_\alpha = R_\alpha \qquad (3)$$

where we have used $P_\alpha = \delta_\alpha n_\alpha T_\alpha$, with $\delta_f = \delta_e = 3$, $\delta_i = 1$, and $S_\alpha^2 = \delta_\alpha T_\alpha/m_\alpha$ and where $R_\alpha$ contains the nonlinear second order terms [9]. Using this set of equations we can now examine the various beam instabilities that can occur in the plasma.

We start by investigating the source of the plasmon field. This can be done by assuming $R_\alpha = 0$, and by neglecting the ion contribution to the high frequency oscillations. We assume perturbations of the form $\tilde{n}_\alpha \sim \exp(i\vec{k} \cdot \vec{r} - i\omega t)$. We can also consider that $\gamma_{0e} \simeq 1$ and $\gamma_{0f} = \gamma_0$. The resulting dispersion relation can be written as $\Omega_e^2\Omega_f^2 = \omega_{pe}^2\omega_{pf}^2/\gamma_0^3$, where we have used the auxiliary quantities $\Omega_\alpha^2 = -(\omega - \vec{k}\cdot\vec{u}_\alpha)[(\omega - \vec{k}\cdot\vec{u}_\alpha) + i\nu_\alpha] + \omega_{p\alpha}^2(1 + k^2\lambda_\alpha^2)/\gamma_{0\alpha}^3$. Neglecting the thermal effects ($S_\alpha^2 = 0$) and collisions ($\nu_\alpha = 0$), we can state the dispersion relation in a much simpler form

$$1 - \frac{\omega_{pe}^2}{(\omega + g\vec{k}\cdot\vec{u}_0)^2} - \frac{\omega_{pf}^2}{\gamma_0^3(\omega - \vec{k}\cdot\vec{u}_0)^2} = 0 \qquad (4)$$

Here we have used for the mean velocity of the return current the value $\vec{u}_e = -g\vec{u}_0$, where $\vec{u}_0 \equiv \vec{u}_f$ is the fast electron beam velocity and the factor $g = n_{0f}/n_{0e}$ is smaller than 1. This equation can be found in Ref. [10], written in a different notation. Assuming that the oscillating frequency $\omega = \vec{k}\cdot\vec{u}_0 + \eta$, where $\eta$ is much smaller than the first term, and considering the quasiresonant conditions such that $\vec{k}\cdot\vec{u}_0 \simeq \omega_{pe}$, we obtain an expression for the maximum growth rate of the electron two-stream instability, which can be written as

$$\Gamma = \frac{\sqrt{3}}{2^{4/3}}\frac{g^{1/3}}{\gamma_0}\omega_{pe} \qquad (5)$$

with $\eta = i\Gamma$. See Fig. 1 for illustration. If we are considering overdense plasma, where $\omega_{pe}^2 \gg \omega_{\mathrm{laser}}^2$, the value of this growth rate can be very significant, specially for ultra intense laser beams in the petawatt domain where the factor $g^{1/3}$ can approach one, even if these high energies imply that the fast electron beam is also highly relativistic, $\gamma_0 \gg 1$ [8,10]. This instability will create electron plasma waves (or plasmon turbulence) with relativistic phase velocities, such that $\omega \sim ku_0 \simeq kc$. This means that the associated plasmons will have a very low group velocity $v_k = S_e^2 k/\omega \simeq S_e^2/c \ll S_e$. Therefore, they are well suited to decay into ion-acoustic waves, as shown below.

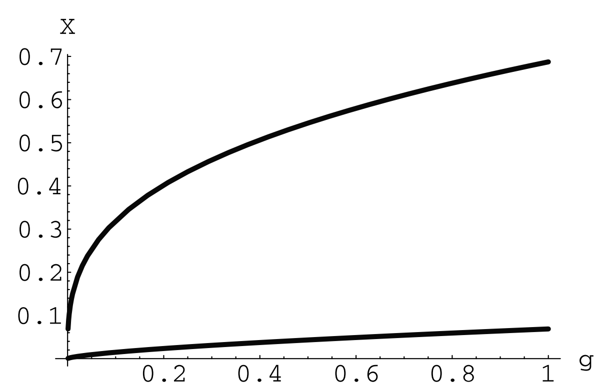


FIG. 1. Normalized growth rate $X = \Gamma/\omega_{pe}$, of the electron two-stream instability, as a function of the fraction of fast beam electrons $g$, for two values of the fast beam gamma factor, $\gamma_0 = 1$ (upper curve) and $\gamma_0 = 10$.

But, before considering the plasmon beam instabilities, we first analyze the possible direct excitation of ion oscillations by the electron beams themselves, the fast beam and the return current drifting across the background ions. Going back to the general Eq. (3) with $R_\alpha = 0$, we get for harmonic perturbations of the type $\exp(i\vec{k}\cdot\vec{r} - i\omega t)$ supported by the return current and the plasma ions, the following dispersion relation, $1 - (\omega_{pi}^2/\omega^2) - \omega_{pe}^2/[\gamma_e^2(\omega - \vec{k}\cdot\vec{u}_e)^2] = 0$, valid in the limit $S_\alpha^2 \simeq 0$, $\nu_\alpha \simeq 0$. In the case of the return current, we have $\vec{u}_e = -g\vec{u}_f$, with $u_f \sim c$, and we can use $\gamma_e = 1$. For moderate values of $|u_e|$ it has been shown that the ion-acoustic waves become unstable, and it has been claimed that this could be a mechanism for ion heating in laser targets as well as in space plasmas [11,12]. Here, however, we have in general a large drift velocity $|u_e| > S_e, S_i$. It has been shown that the electron-ion oscillations also become unstable [13], with a maximum growth rate of $\Gamma_{\max} = \sqrt{3}(Zm_e/2^4 m_i)^{1/3}\omega_{pe}$ corresponding to a frequency of the most unstable modes of $\omega_{\max} \simeq \Gamma_{\max}/\sqrt{3}$ and a wave number $k_{\max} \simeq \omega_{pe}/u_e$, but this frequency is very large as compared with that of the ion-acoustic oscillations, which means that the oscillations would be mainly supported by the electrons and not the ions. On the other hand, the phase velocity of these waves would be much larger than the ion sound speed. So, these electron-ion oscillations are not a plausible support for the energy transfer to the ions, and will be disregarded.

In order to study the influence of the plasmon beam on the ion wave stability we go back to Eq. (3) and retain the nonlinear terms $R_\alpha$. Neglecting the contribution of the fast electron beam and assuming a sinusoidal perturbation, we can derive a nonlinear dispersion relation for low frequency oscillations in the presence of the plasmon turbulence of the form

$$\left(\Omega_e^2 - \omega_{pe}^2\frac{\omega_{pi}^2}{\Omega_i^2}\right) = Q_e + ZQ_i\frac{\omega_{pi}^2}{\Omega_i^2} \qquad (6)$$

where $Q_\alpha$, with $\alpha = e, i$, are very complicated expressions depending on the equilibrium plasmon distribution $N_0(\vec{k}')$, to be specified later. In the absence of plasma turbulence, we have $Q_e = Q_i = 0$, and this equation reduces to $\Omega_e^2\Omega_i^2 = \omega_{pe}^2\omega_{pi}^2$, which in the absence of collisions and in the low frequency range are reduced to the well-known dispersion relation for ion-acoustic waves $\omega^2 = k^2 v_s^2/(1 + k^2\lambda_e^2)$, where $v_s = \sqrt{(3ZT_e + T_i)/m_i}$ is the ion sound speed. In the presence of plasma turbulence, these waves can become unstable. Let us assume the simple case of a plasmon beam described by the unperturbed plasmon distribution $N_0(\vec{k}') = (2\pi)^2 N_0 \delta(\vec{k}' - \vec{k}_0)$. In this case, we get, for parallel propagation,

$$Q_\alpha = \frac{\hbar N_0}{n_o m_\alpha}\frac{\omega_{p\alpha}^2}{\omega_0'^3}\frac{\omega_{pe}^2 k^2 S_e^2}{[(\omega - ku_0) - 2i\Gamma_0]^2}. \qquad (7)$$

Replacing this in Eq. (6) leads to the following dispersion relation, valid for ion-acoustic waves in the presence of a plasmon beam

$$\Omega_e^2\Omega_i^2 = \omega_{pe}^2\omega_{pi}^2\left\{1 + \frac{\Omega^2}{[(\omega - ku_0) - 2i\Gamma_0]^2}\right\} \qquad (8)$$

where $\Gamma_0$ is here the growth rate of the unstable plasmon spectrum and $\Omega^2 = (ZW_0/n_0 m_e c^2)(k^4 S_e^2 c^2/\omega_0'^2)$. It is clear that $\Omega$ has the dimensions of a frequency, which can be seen as an effective plasmon beam frequency. We have used here the energy density of the plasmon beam, as defined by $W_0 = \hbar\omega_0' N_0$. This expression generalizes our previous result [9] to the case of a relativistic and collisional plasma. It can easily be seen from such a dispersion relation that the ion-acoustic waves become unstable. In order to determine the maximum growth rates, let us consider the case of a nearly resonant situation where the phase velocity of the low frequency wave is nearly equal to the plasmon beam velocity.

If we go back to the definition of $\Omega_\alpha$ we see that, for ion-acoustic waves, we can use the approximation $\Omega_e^2 \simeq \omega_{pe}^2(1 + k^2\lambda_e^2)$. In order to study the stability of these waves in the presence of the plasmon beam and to derive the expression for the maximum growth rates, we can also assume that $\omega = \omega_s + \eta$, with $\omega_s = kv_s(1 + k^2\lambda_e^2)^{-1/2}$. The above dispersion relation can then be reduced, for $\omega_s \simeq ku_0$ and $\eta \ll \omega_s$, to the simple expression: $(2\eta + i\nu_i)(\eta - i2\Gamma_0)^2 = \Omega^2\omega_s$. In the absence of dissipation, this would simple lead to $\eta = i\Gamma$ where the growth rate of the ion-acoustic waves would be given by [9]: $\Gamma = \sqrt{3}(\omega_s\Omega^2/2^4)^{1/3}$. This upper limit of the growth rate is illustrated in Fig. 2. Notice that the quantity $Y_s = \Omega/\omega_s$ can be approximately written as $Y_s = (Z\gamma_0 g)^{1/2} \times (\omega_s/\omega_{pe})(c/S_e)$, where we have assumed that the plasmon energy is of the order of the fast electron beam energy. It is obvious that, for large $\gamma_0$ and low electron temperature, this quantity can be in the range of 0.1–1, thus leading to very high growth rates comparable to the those of the

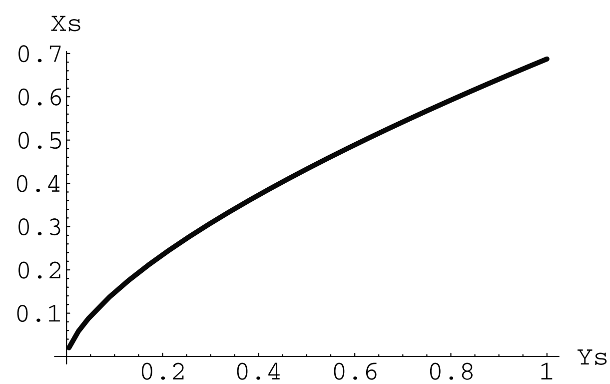


FIG. 2. Normalized growth rate $Xs = \gamma/\omega_{pe}$, of the ion-acoustic waves, as a function of the normalized plasmon intensity $Ys = \Omega/\omega_s$, for $\Gamma_0 = \nu_i/4$.

electron two-stream instability, which is in the range of $10^{-1}$–$10^{-2}\omega_{pe}$.

For a dissipative and unstable system, we obtain the approximate expression

$$\Gamma = \frac{\sqrt{3}}{2^{4/3}}(\omega_s \Omega^2)^{1/3} + \frac{4}{3}\Gamma_0 - \frac{\nu_i}{6}. \tag{9}$$

We see that, for a moderate level of the electron two-stream instability, such that $\Gamma_0 > \nu_i/8$, the ion-acoustic waves will always become unstable, even for a very small value of the plasmon turbulence. The electron two-stream instability will then be saturated by the fast energy transfer to the ion-acoustic wave spectrum, which in turn will be heavily damped due to the high rate of ion collisions. Such a two-step process will then efficiently transfer energy from the fast electron beam (in reality, from the incident laser beam) to the ion thermal energy without heating the electrons. The electrons only mediate the transfer, without being significantly heated, because they have high mean velocities (which corresponds to low collision frequencies) and support high frequency plasma wave oscillations which are not significantly affected by collisions. In the end, these physical processes could result in a scenario where the bulk ion heating is greater than the electron heating if the ion growth rate exceeds the rate at which the energy in the electron oscillations is eventually converted to electron heating. The condition for preferred ion heating to occur will then be the formation of a sufficiently intense fast electron beam such that $g > 2^4(\gamma_0 \nu_e/\sqrt{3}\omega_{pe})^3$, where $g = n_{0f}/n_{0e}$. This condition can also be cast in terms of the laser intensity and the plasma density and temperature.

Assuming near net balance between the fast electron current and the return current, we can write $u_e = n_{0f}u_f/n_{0e}$. On the other hand, the energy of the fast electron beam will be a fraction of the laser beam energy, which can be stated as $n_{0f}K_f u_f = f_{\rm abs} I$, where $K_f$ is the mean kinetic energy of the fast electrons and $f_{\rm abs}$ is the laser absorption factor into fast electrons. Using $u_f \simeq c$ and $K_f \simeq \gamma_0 m_e c^2$, we can then obtain $u_e \simeq f_{\rm abs} I/\gamma_0 n_{0e} m_e c^2$. Substituting this into the expression for $g$, we get a threshold criterion expressed in terms of the laser beam intensity, as

$$I > \frac{2^4}{3^{3/2}} \frac{n_{0e} m_e c^3}{f_{\rm abs}} \gamma_0^4 \left(\frac{\nu_e}{\omega_{pe}}\right)^3 \tag{10}$$

Notice that the relativistic gamma factor $\gamma_0$ and the electron return velocity $u_{0e}$ are typically proportional to $\sqrt{I}$, and that the collision frequency goes with $u_{0e}^{-3}$. This means that the right hand side of this inequality varies with the laser intensity as $I^{-5/4}$. For a typical laser target experiment with $n_{0e} \sim 10^{23}$ cm$^{-3}$, this threshold criterion can only be satisfied for laser intensities in excess of $10^{20}$ W cm$^{-2}$. On the other hand, the collision frequency $\nu_e$ is also proportional to $Zn_{0e}$. This then provides an explanation for the recent high conversion efficiency in the integrated cone-guided experiments in Japan [1,2]. The strong $Z$ and density dependence of Eq. (10) makes it more difficult to excite ion heating in the gold cone-guide and the energy transfer is greatest in the compressed carbon-deuterium plasma rather than in the gold cone guide (where the energy loss is Ohmic).

Equation (10) also explains why the heating is localized between the cone tip and the compressed core in the simulations of Campbell *et al.* [4]. The coupled instability and its associated ion heating is generated only in these lower density regions of the target when the intensity is below $10^{20}$ W cm$^{-2}$. The usual equilibration process will then tend to equalize the ion and electron temperatures below 1 KeV. The nonlinear saturation regime of this two-step process will be discussed in a future work.

In conclusion, we have proposed a new mechanism for anomalous ion heating in ultraintense laser plasmas. This mechanism is based on the excitation of two coupled instability processes that drive the laser energy down to the ion population. First, a strong electron two-stream instability is driven by the fast electron beam generated by the incident laser beam. Second, the resulting high phase velocity plasma waves decay by a resonant coupling process into ion-acoustic oscillations. These low frequency ion oscillations are then strongly damped by the ion collisions in the dense plasma. The present theoretical model gives a simple explanation for the preferential heating of the bulk ion population observed in various laser experiments recently performed in the petawatt regime. Such an anomalous heating is of paramount importance in the fast ignition scenario for inertial fusion.

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