Evidence of anomalous resistivity for hot electron propagation through a dense fusion core in fast ignition experiments

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Abstract. Anomalous resistivity for hot electrons passing through a dense core plasma is studied for fast ignition laser fusion. The hot electrons generated via the ultra-intense laser pulse and guiding cone interactions are measured after they pass through a dense plasma with a density of 50–100 g cm\(^{-3}\) in a radius of 15–25 µm. When significant neutron enhancements are achieved by the ultra-intense laser pulse injection, the energy reduction of fast electrons is observed. Also, a reduction in the number of electrons with energy up to 15 MeV can be seen. We offer a new physical mechanism for the stopping of electrons, involving electron magnetohydrodynamic shock formation in the inhomogeneous plasma density region. The dissipation in the shock region can explain electron stopping with energies of the order of 15 MeV.

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1. Introduction

Human endeavors to harness energy by nuclear fusion have led to enormous progress in science and technology in the last four decades, but the dream of power production has remained ever elusive. Efforts to realize laser fusion received a major boost with the fast ignition proposal in 1994 [1]. This proposal seeks to implode/compress a fusion pellet by multiple, high-energy, nanosecond laser pulses and then strike a spark in the compressed core at the time of maximum compression. The spark is provided by mega-amperes (MA) of a hot (MeV) electron current created by a separate, ultra-short pulse laser. The attractiveness of the FI scheme in terms of (a) the separation of the compression and ignition steps and (b) the relaxation of the symmetry otherwise required in the implosion process has stimulated investigations by laboratories around the world. In 2001, the first support for its feasibility was offered in experiments where neutron emission was seen to be enhanced by a factor of 1000 when a deuterated-polystyrene (CD) pellet compressed to 55 g cm$^{-3}$ density was ignited by a few MeV and a few tens of MA of hot electron current generated by a petawatt laser pulse [2, 3]. Hopes thus engendered of the feasibility of FI, however, continue to encounter skepticism regarding the appropriate transport of MA electron currents through the tens of microns distance of the compressed core at number densities as high as $10^{25}$ cm$^{-3}$. This skepticism is fueled by (a) the considerable existing knowledge of the complexities of relativistic electron transport at normal densities on the one hand and (b) the total paucity of experimental information of transport through high-density cores on the other [4]–[7]. Here we highlight this deficiency by offering direct, quantitative experimental evidence of the fast ignition process. Firstly, we compare the number of electrons traversing the dense core with that passing through the low-density coronal plasma. We observe enhanced stopping of the electrons through the dense core in the event of successful (higher neutron yield) fusion reaction. Secondly, we offer a physical mechanism for the stopping of electrons by invoking electron magneto-hydrodynamic (EMHD) processes for the electron transport. The mechanism is based on the organization of the energetic electron current in the form of sharp current layers in the presence of a strong background plasma density gradient. The energy dissipation is typically high in such sharp current layer patterns. This is because the sheared current layers are unstable to the excitation of Kelvin Helmholtz instabilities [8]–[10] and the resulting electromagnetic turbulence leads to energy cascades to short scales and eventual dissipation through Landau damping, a process that may be modeled by anomalous enhancement of viscosity coefficients for the energetic electron fluid. This particular mechanism is especially important in the present context as the other often cited mechanisms of stopping, like the Weibel and electrostatic instabilities, are either inoperative or have the wrong scaling with density and hence are irrelevant in the dense core.
2. Hot electron observation in integrated experiments

The FI scheme has been demonstrated using a guiding gold cone at the Gekko XII and PW laser facility, Osaka University as reported by Kodama et al [2, 3]. The compression of a spherical shell with a gold cone is achieved with nine beams of GXII laser with a total energy of 2.5 kJ of the wavelength $\lambda = 532$ nm in a pulse duration of 1.2 ns. Here, the shell is made with CD with a radius (thickness) of $250 \mu$m ($7 \mu$m). The cone is made of gold and its tip is capped by a $5 \mu$m thick gold foil and has a $30^\circ$ opening angle. The profile of imploded core plasma is estimated from measurement with an x-ray backlight method. A high dense core with a density (radius) of $50–100$ g cm$^{-3}$ ($15–25 \mu$m) is created at the front of the cone by high energy, nanosecond implosion laser pulses [2, 3]. A PW class ultra-intense laser (UIL) pulse ($\lambda = 1 \mu$m) [11] is injected into the cone with an energy of $200–300$ J in a pulse duration of 0.5 ps. In the experiment, neutron yield is measured with the time-of-flight methods.

Two electron spectrometers (ESMs) measure the energy spectra of hot electrons in vacuum at two different positions. One ESM is arranged at $20^\circ$ from the UIL propagating axis and the other is set at $40^\circ$ as shown in figure 1. The ESM $20^\circ$ views electrons passing straight through the dense core, while the ESM $40^\circ$ measures electrons that are essentially passing through the low-density coronal region of the plasma. The solid angle of ESMs is limited as $1 \times 10^{-5}$ Sr by a collimator with a diameter (length) of 5 mm (40 mm). The electrons are bent by the magnetic field (up to 0.45 T) after they pass through the collimator and then enter a detector, imaging plate (IP). The sensitivity of the IP for electrons has been absolutely calibrated [12]. The ESMs have an energy resolution ($E/\Delta E$) over 40.

When the UIL pulse is injected into the cone at the time of maximum compression, we observe a reduction in transmitted electron flux compared to the mistimed injection, with good reproducibility. This is accompanied by increased neutron yield, providing direct indication that the energy of the hot electrons is responsible for the fusion in the core. Figure 2 presents the electron energy spectra measured after it has transited through the plasma. The injection timing shown in the figure is measured from the timing when the maximum neutron yield is achieved. Figure 2(a) shows the electron flux measured directly behind the core. The solid and dotted lines represent fluxes for the cases where the yield enhancement of thermal neutrons, i.e. increased fusion, is observed. When the UIL pulse is injected at $+50$ ps or $-170$ ps, significant neutron

Figure 1. The UIL pulse is injected into the gold cone with an imploded plasma core with a density of $50–100$ g cm$^{-3}$. Two ESMs measure electrons coming out at $20^\circ$ and $40^\circ$ from the laser axis. Electrons received by the ESM $20^\circ$ pass through the dense core. The ESM $40^\circ$ measures electrons passing through the corona plasma.
The electron energy reduction observed with neutron enhancement strongly indicates that the core is heated by the hot electrons. The core temperature increase is estimated with the observed energy reduction of hot electrons. The electron angular distribution is estimated from the experimental measurements with ESMs. Then, the number reduction of hot electrons indicates that the energy of \(2 \times 10^{14}\) MeV is deposited in the core when a 40% conversion enhancement is not observed. The crucial observation is that the total energy of electrons measured on the spectrometer is reduced by as much as 44%. The fluxes measured with the ESM 40° shown in figure 2(b) indicate insignificant modification of the transport of hot electrons in the low-density plasma corona. Here, it is worth describing the shot-to-shot fluctuations on the electron energy spectra. The electron number reduction is observed repeatedly at around the maximum compression as shown in figure 2(a). The reduction of 44% is beyond the fluctuations comparing against the similarity of the shots at the mistiming, i.e. +50 ps and -170 ps. The electron number and the feature of energy spectra could be varied if the tip of the cone is broken before the UIL pulse interacts with the cone. However, the spectra are, again, similar before and after the maximum compression at both 20° and 40°. These results indicate that the laser and cone target interactions are not changed with the UIL injection timing.

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efficiency from the laser to the hot electrons is assumed. The deposited energy explains the making of a 15 µm hot spot (in radius) with a temperature of 1.1 keV in the core with a density of 55 g cm\(^{-3}\) after a temperature increase by 700 eV due to the hot electrons as reported in [3]. The temperature is consistent with the experimental results estimated from the neutron measurements. It is interesting to see that electrons as energetic as 15 MeV are stopped by the core as shown in figure 2(a), in contradiction to the prevalent notion that only low-energy electrons (up to 1 MeV) can be stopped by a dense core [4]–[7].

3. Theoretical mechanism for hot electron stopping through a dense core

Our experimental results clearly show that electrons with energies as large as 15 MeV propagating through the dense target core region get stopped. However, electrons passing through the low-density coronal region are relatively unhindered. Here, we identify a mechanism that explains these observations. As energetic electrons from the critical density surface of the UIL move towards the target core, a return shielding current is induced by the plasma. This return current helps in neutralizing the space charge and the magnetic fields generated by the electron beam current. It is well known that forward and return shielding currents become spatially separated via Weibel instability in the low-density plasma corona region \(n_0 \sim n_b \sim 10^{21}\), forming cylindrical current channels propagating towards the target core. The center of the cylinder carries the forward current, which is surrounded by a cylindrical shell of return plasma current. The crucial question concerns the mechanism of energy dissipation of fast electrons within this current channel as it propagates towards the strongly inhomogeneous background plasma associated with the compressed core. The mechanism clearly has to rely on processes other than direct collisional stopping of fast electrons since that process is relatively weak for parameters of the experiment. The mechanisms proposed so far have relied on electrostatic instabilities [13], Weibel instability and coalescence of current filaments. These processes, however, are important only in the low-density coronal region as their growth rates/effectiveness scale adversely with the ratio of fast to background electron densities.

A possible way by which the current pulse, associated with energetic electrons going inwards, can dissipate its energy in the high-density plasma of the target core region is if the current layers can be sharpened by some mechanism. Note that with the separation of the forward and return currents through Weibel separation significant self-consistent magnetic fields are generated, which influence the motion of all electrons such that a description of current pulse propagation by EMHD theory becomes relevant [5]. We show that in the presence of an inhomogeneous plasma there is indeed an EMHD mechanism by which the current layers can be sharpened, leading to the formation of current shocks. The width of these shock structures adjusts in such a fashion that the energy dissipation is independent of both the magnitude and the nature of the underlying dissipation mechanism of the region. Thus, even when the value of the classical resistivity coefficient is small, the damping is significant due to the sharpening of the layers. In fact the shock width adjusts itself such that the magnitude of energy dissipation is independent of the value of plasma resistivity. In the event that the classical resistivity parameter is extremely small, the shock structure in the presence of plasma density inhomogeneity would keep getting sharper and would hit the collisionless electron skin depth. At this point the electron inertia-driven Kelvin–Helmholtz (KH) instabilities [8]–[10] would be excited, which would lead to turbulence generation and an effective anomalous viscosity. This is a completely collisionless mechanism of energy dissipation of the current channel. The KH instabilities are

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relatively independent of the fast to background electron density ratio \([9, 10]\) and convert the electron flow energy into fine scale vortices. In the three-dimensional (3-D) case such vortex flows cascade energy towards finer scales and eventually dissipate into heat by electron Landau damping effects. The shock width in the presence of anomalous viscosity would adjust and stabilize at a value governed by the anomalous viscous coefficient. It, however, again yields the same magnitude of energy dissipation. Hence for any practical purposes it becomes irrelevant whether the energy dissipation is due to classical resistivity and/or anomalous viscosity.

We now provide a detailed description of the EMHD mechanism of the sharpening of current layers. The generalized EMHD evolution equations (including electron inertia-related terms) for an inhomogeneous density plasma can be written as \([14, 15]\)

\[
\frac{\partial \vec{g}}{\partial t} = \nabla \times \left( \vec{v} \times \vec{g} \right) - \eta \nabla^2 \vec{B},
\]

\[
\vec{v} = -\frac{1}{n} \nabla \times \vec{B}; \quad \vec{g} = \frac{\nabla^2 \vec{B}}{n} - \nabla \left( \frac{1}{n} \right) \times (\nabla \times \vec{B}) - \vec{B}.
\]

The fields \(\vec{v}\), \(\vec{B}\) and \(n\) in equations (1) and (2) denote normalized electron velocity, magnetic field and plasma density, respectively. Equation (1) has been obtained by taking the curl of the electron momentum equation. The first two terms in the expression of \(\vec{g}\) in equation (2) are the electron inertia-dependent terms of the electron momentum equation. The ions do not respond at fast EMHD time scales and merely provide a static neutralizing background. The displacement current contribution is ignored under the EMHD approximation. This relates the electron velocity to the magnetic field directly through Ampere’s law in equation (2). The collisional friction between electron and ions has been retained through the resistivity coefficient \(\eta\).

We have chosen to normalize the electron density with a typical value of the plasma density \(n_{00}\) in the region of interest, and the length by the corresponding skin depth \(d_{e0} = c/\omega_{pe0}\) (\(\omega_{pe0} = 4\pi n_{00}e^2/m_e\)) corresponding to this density. The magnetic field is normalized by \(B_0\) (the typical magnitude of the magnetic field) and the time has been normalized by the corresponding electron gyrofrequency \(\omega_{ce0} = eB_0/m_ec\). It should be noted that equations (1) and (2) involve no approximations other than the stationarity of ions and the neglect of the displacement current.

It was shown analytically by Kingsep et al. \([16]\) that the propagation of EMHD structures through an inhomogeneous plasma leads to current pulse sharpening and shock formation. Although the analytic derivation was restricted to an inertialess EMHD fluid, Kingsep had also pointed out that the same phenomena would take place even when finite electron inertia-related effects are included. This has been explicitly shown recently by us \([14, 15]\) by numerically simulating equations (1) and (2) (which contain electron inertia terms), the propagation of a current pulse through an inhomogeneous plasma. The simulations also demonstrate that the width of the shock region is not constrained by the local skin depth of the plasma medium. The amount of energy dissipated through this shock region was also found to be independent of the magnitude as well as the character of underlying dissipation in the simulation.

The physical reason for shock formation can be understood as follows. The current pulse structure moves with an axial drift velocity due to electron inertia-related terms. However, in the presence of an inhomogeneity an additional drift velocity \(\vec{V}_n = \vec{B} \times \nabla n/n^2\) for the structure exists. This can be observed by substituting for \(\vec{g}\) from equation (2) to equation (1). Here \(\vec{B}\) is the magnetic field associated with the structure. For a cylindrical current pulse structure \(\vec{B}\) is along the \(\hat{\theta}\) direction and the density inhomogeneity is along \(\hat{z}\). This drift is radially inward for a
current pulse moving towards denser plasma region, and is thus responsible for radial pinching of the structure and shock formation. The simulations of references [14, 15] were carried out for a slab representation and clearly showed the pinching of the current pulse transverse to both the magnetic field and the inhomogeneity. Another explanation for the shock formation can be given by realizing that the $-\vec{B}$ term of $\vec{g}$ produces an additional nonlinear source term in equation (1) (in the presence of inhomogeneity). This nonlinear term has the form of $\nabla n \times \nabla B^2/2$ in the evolution equation. Such a term is similar to the well-known ‘thermoelectric’ current drive proportional to $\nabla n \times \nabla p$ ($p$ being the thermal pressure) thoroughly studied in ideal MHD fluid equations. Here, however, the magnetic pressure itself plays the role of $p$, making the source term nonlinear (e.g. dependent on $B^2$). The nonlinearity of the source produces steepening and sharpens the current layers.

We now apply equations (1) and (2) to illustrate EMHD shock formation. Under a slab representation with $\hat{x}$ for the radial and $\hat{y}$ for the azimuthal direction, and with the approximate expression $\vec{g} \approx -b\hat{\theta} = -b\hat{y}$, equation (1) in the inertialess limit reduces to the following dissipative Burger’s equation [15]:

$$\frac{\partial b}{\partial t} + Kb \frac{\partial b}{\partial x} = \eta \frac{\partial^2 b}{\partial x^2}. \tag{3}$$

Here, $K = d_e/\alpha_n$, with $\alpha_n = n/\sqrt{\nabla n}$ as the density gradient scale length. The Burger equation is well known for supporting shock solutions. Thus when a propagating current filament encounters a sharp density gradient it steepens in the radial direction forming radially propagating shocks. The shock width gets determined by the balance of the nonlinear term with the classical resistivity parameter. This suggests that the radial shock width $l_x$ would scale as $l_x \sim l_w = \eta/Kb$. The net energy dissipation over a length $L$ in this case is

$$Q = \int_a^b \int^L \int_{l_x}^b \eta \left( \frac{\partial b}{\partial x} \right)^2 dx \, dz \, dy \tag{4}$$

$$= \frac{b^2 a L l_x}{l_x^2} = b^3 K L a = b^2 a^2 K L v_e.$$  

Here $b \sim a v_e$ ($v_e$ being the electron velocity) from Ampere’s law has been used. Here the shock length $z = L$, the shock width $l_x \sim l_w = \eta/Kb$ and $a$ denotes the transverse extent. Note that the energy dissipation is independent of the magnitude of the resistivity. As stated earlier, this arises because the thickness of the shock and the current spike adjust in such a manner as to cancel the resistivity coefficient.

Although the above derivation is restricted to inertialess EMHD fluid, we have recently verified by numerical simulations that similar effects take place in the presence of inertia-related terms [15]. In the presence of inertia another physical effect is also important. When the value of the classical resistivity parameter is extremely small, it is likely that $l_w$ would be much smaller than electron skin depth. In this case, while the shear width of the current layer approaches the value of $l_w$, it would first hit the electron skin depth. However, as soon as the shear in the current layer becomes comparable to the electron skin depth the current layer becomes susceptible to KH instability [8, 9] due to inertial effects. The ensuing turbulence in 3-D leads to an effective anomalous viscous dissipation [10]. Incorporating the turbulence-induced effective viscosity $\mu_a$ in equation (3), we obtain

$$\frac{\partial b}{\partial t} + Kb \frac{\partial b}{\partial x} = \eta \frac{\partial^2 b}{\partial x^2} - \mu_a \frac{\partial^4 b}{\partial x^4}. \tag{5}$$

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In this case \( \mu_a \) dominates over classical resistivity and the shear width gets decided by the balance of the nonlinear term with viscous dissipation due to the anomalous term. The radial shock thickness in this case scales as \((\mu_a/Kb)^{1/3}\) and the energy dissipation as shown below is again independent of the magnitude of \( \mu_a \). A net energy dissipation rate \( Q \) in a length \( L \) scales as \( \sim f(\mu_a(b^2 d^2/dx^2) d^2x) 2\pi aL \sim \mu_a 2\pi aK L b^2 / x^3 \sim 2\pi aK L b^3 \). Again, using \( b \sim av_e \) gives \( Q \sim K L b^3 a^2 v_e \). Thus, the evaluation of energy dissipation is not dependent on any approximations that would be necessary for estimating the value of \( \mu_a \).

Another interesting feature of the estimated energy dissipation is the fact that when \( KL \sim 1 \), the dissipation equals \( b^2 a^2 v_e \), which is the total magnetic energy influx rate in the inhomogeneous region due to the propagating current channel. In a Weibel separated current channel, the ratio of kinetic to magnetic energy gets decided by the ratio of its width compared to the skin depth \( (mnv_e^2)/(B^2/8\pi) = (c^2 m)/(2\pi ne^2 a^2) = (2c^2)/(\omega_{pe}^2 a^2) \). Here \( B = (4\pi/c) a n e v_e \) has been estimated from \( \nabla \times \vec{B} = (4\pi/c) \vec{J} \). Thus if a Weibel separated current channel (coming towards the inhomogeneous plasma region) has an extent that is comparable or wider than the electron skin depth, the magnetic energy constitutes a significant portion of its total energy. In this case then, \( L \sim K^{-1} \sim L_n \) (the background density gradient scale length) essentially matches the stopping length \( L_s \) for the electrons independent of the detailed nature of microscopic dissipation in the sharpened current layer.

Thus, the sharpening of current layers in the region where the plasma density gradient is high plays a crucial role in the dissipation of energy associated with energetic electrons. It should be emphasized here that results published from a number of PIC simulations ([6] and references therein), [17]–[19] on the propagation of energetic electrons through a background inhomogeneous plasma show evidence of efficient heating of the background ions at the spatial location where the background plasma gradient is present.

4. Discussion and summary

In the experiment, the hot electron beam measured at 40° travels through the low-density plasma corona avoiding the sharp density gradient region and is not stopped. On the other hand, the 20° beam that passes through the dense core at the correct UIL injection timing encounters the sharp density gradient region and hence gets stopped within \( L_n \sim 20 \mu m \). The resistance \( R \) offered by the compressed target core may be written as \( R \sim Q/I^2 \), where \( Q \) is the energy dissipation rate and \( I \) is the current carried by the filament. Taking \( Q \sim (B^2/4\pi) a c^2 \) and \( B \sim 2I/ac \), we get, \( R \sim 1/c \sim 30 \Omega \). Thus, for a typical current filament carrying 0.5 MA of hot electron current, electrons up to an energy of order \( 1R \sim 15 \text{ MeV} \) may be stopped within the stopping length \( \sim 20 \mu m \), in agreement with the experimental observations.

In conclusion, we have provided the first, direct and semi-quantitative proof of energy dissipation of fast electrons in the compressed core of a scaled-down fast ignition experiment. The electron energy spectra showed that electrons up to 15 MeV appear to be stopped losing 44% of their energy. This proof is provided by the direct, well-calibrated measurement of MeV electron stopping by the core. We also offer a clear, well-formulated physical mechanism that explains not only our observations but also other simulation studies, thereby opening a new perspective on this problem. We believe that our results will provide the confidence necessary to move forward in the realization of fast ignition on an economically and technologically viable scale.
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