Intense laser-generated ion beams propagating in plasmas

A P L Robinson

Central Laser Facility, STFC Rutherford-Appleton Laboratory, OX11 0QX Didcot, United Kingdom

E-mail: alex.robinson@stfc.ac.uk

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Abstract
The generation of nuclear reactions using laser-generated proton and ion beams has been investigated for several years by a number of different research teams. Interpreting and understanding these experiments remains a challenge, which suggests that there are phenomena, processes, or effects not considered in the interpretive framework that is usually employed. In this paper, we examine the theoretical case for strong drag heating by the incident ion beam, and the generation of electric fields, which affect the propagation of the ion beam and thus the yield from nuclear reactions. We obtain an analytic metric, $\Xi_h$, for predicting the regime in which the drag heating heats the target so violently that ion–electron drag is suppressed. This compares favourably to ion kinetic hybrid simulations.

Keywords: laser, ion beam, target heating

1. Introduction
In recent years, the study of nuclear reactions driven by ultra-intense laser irradiation has been highly active in many different research groups and institutes [1–12]. In the majority of cases these nuclear reactions are thought to occur due to the generation of energetic proton or ion beams (a well-established phenomena in laser–solid interactions [13–20]) which then propagate into the irradiated target, a secondary target, or other surrounding material.

Interpreting and understanding the results of these experiments appears to be more challenging than might be expected. In a number of experiments there has been a recurring theme where experimental results are at odds with theoretical models. While one can conjecture circumstances in which experiment and theory are reconciled, how this is done self-consistently is currently not clear.

For example, in a recent paper by Giuffrida et al [4], which was concerned with proton–boron reactions in laser-irradiated solid targets, the observed high alpha yield could be reasonably produced provided a sufficiently thick layer of hot (1 keV) material was present. The issue is that it is not clear that this is consistent with radiation-hydrodynamics simulations, or any other modelling of the experiment. Another example of this can be found in the significantly earlier report by Toupin et al [2], which attempted to model neutron production in experiments involving deuterated targets. This, and other reports in the current literature, points to problems in our understanding of ion beam transport in dense plasmas, at least in certain parameter regimes. Interestingly, these issues suggest that the range of energetic ions is much greater than expected, but in order to resolve this problem one would have to show that there is an effect that could increase the ion range by at least a factor of 10.

In this paper, we re-examine the standard approaches to ion beam-target interactions in laser–solid studies. We show that the heating of the target by ion–electron drag can be so strong that the ion–electron drag is substantially reduced, and the penetration range of the ion beam is substantially enhanced. We derive an analytic measure ($\Xi_h$) for when this regime is
entered, and show that this measure effectively predicts the onset of this regime in numerical simulations using an ion kinetic hybrid code. In our numerical results we observe that the ion range can increase by of 11.3, which indicates that the drag heating due to the incident ion beam itself could play a major role in resolving the outstanding problem that we highlight.

2. Notation

As we are primarily concerned with a high energy ion or plasma beam that is incident on a static target plasma, direct reference to the target ions or ‘background’ ions is seldom required. Thus, when we refer to electron quantities, e.g. $n_e$, this should be taken as a reference to the target or background (i.e. $n_e$ is the electron density of the target plasma). Likewise, when we refer to ion quantities, this should be taken as a reference to the incident ion or plasma beam. Such terms, such as $v_i$, $Z_i$, and $n_i$ refer to the ion velocity, ion charge state, and ion density of the incident ion beam and not that of the target/background plasma. We use this convention throughout this paper, unless stated otherwise. Note that we also work in units of eV for temperature.

3. Current treatment of ion transport

The standard approach to ion beam transport in many laser–solid studies can be summarized as follows:

(a) Single-particle treatment of the ions.
(b) Ions are subject to drag and angular scattering only.
(c) Target conditions are assumed to be static, and usually homogeneous.
(d) Ion transport does not alter the target conditions.

Examples of this include those in [2, 4], but also in [21]. For the purposes of numerical simulations this approach lends itself very well to Monte Carlo methods [21]. In many cases, a full justification for this approach is not given. This approach can only ever be an approximation and can be a very good approximation, but on its own, it is certainly not fully self-consistent. In the regime where the ion beam is very dilute in comparison to the target density, then we can reasonably expect this ‘single particle’ framework to be a very good approximation. In the case of laser-generated ion beams there are a variety of scenarios, and, in some of these, it is possible that the ion beam will no longer be highly dilute in comparison to the target density.

What might happen in this case? Since ions are subject to drag, there is energy transfer from the ions to the background material. Under certain conditions the heating could become so strong that it will affect the stopping power (as this is temperature dependent). If the heating is strong, then this implies the existence of significant temperature gradients. In a plasma, this could eventually become so significant that we need to account for the electric and magnetic fields associated with such temperature inhomogeneities.

In addition to the temperature gradients, the ion drag implies an equal but opposite force acting on the electron fluid. In order for the electrons to remain in momentum balance, there should be an additional contribution to the electric field that will be described by an appropriate modified Ohm’s law. Normally this is negligible, but it is conceivable that if the ion beam density is sufficiently large, then the electric field will act to significantly mitigate the ion drag.

These possibilities can be conjectured just on qualitative grounds alone, however one can immediately see that these possibilities are only likely to become important if the ion beam density is no longer a tiny fraction of the target density. Further exploration requires quantitative analysis, which we shall proceed to look at in the next section.

4. Heating of plasma by ion drag

The standard correction to the stopping power in a finite-temperature plasma comes in the form of a multiplying factor that depends on $v_i/v_{th,e}$ ($v_i$ is the ion beam velocity, and $v_{th,e}$ is the electron thermal velocity) [22–24]. That is the temperature independent stopping power,

$$\frac{dE}{dx} = \frac{Z_i^2 e^4 n_e}{4\pi\varepsilon_0 m_i v_i^2} L, \quad (1)$$

(where $L$ is the stopping number), becomes

$$- \frac{dE}{dx} = \frac{Z_i^2 e^4 n_e L}{4\pi\varepsilon_0 m_i v_i^2} G \left( \frac{v_i}{v_{th,e}} \right), \quad (2)$$

where $G$ is the function that corrects the stopping power due to the finite temperature of the electrons. Physically this effect emerges because the electron temperature affects both binary collisions and plasma wave excitation. In particular, once the electron thermal velocity approaches the ion beam velocity, one expects the stopping power to be significantly curtailed. The correction that is often applied, the Chandrasekhar correction, has the form,

$$G(y) = \text{Erf} \left( \frac{y}{\sqrt{2}} \right) - \sqrt{\frac{2}{\pi}} ye^{-y^2/2}, \quad (3)$$

where $y = v_i/v_{th,e}$. For the purposes of this study, the level of accuracy that this affords is sufficient, although it can be improved upon (see [22]). From this we can indeed see that $G(1) \sim 0.2$. In terms of gauging the significance of heating, we can therefore take $v_{e,th} = v_i$ to determine a ‘critical temperature’,

$$T_{crit} = \frac{m_e v_i^2}{3e}. \quad (4)$$

The temperature evolution due to ion stopping alone will be,

$$\frac{3}{2} e n_e \frac{\partial T_e}{\partial t} = f_{d\text{rag}} v_i n_i, \quad (5)$$
for a monoenergetic ion beam (where \( f_{\text{drag}} = -\frac{dE}{dx} \)). For a beam with significant energy spread the RHS becomes an integral over the distribution function. For the purposes of exploring the issue of the drag heating the monoenergetic beam model will suffice (the integral over the distribution function is incorporated in the later numerical model of section 8). From this we can approximately determine the conditions under which the heating via drag will be sufficient to reach \( T_{\text{crit}} \) during the ion beam pulse duration, \( \tau_i \), via,

\[
\frac{f_{\text{drag}}^* v_i \mu_i}{3e_n} > \frac{T_{\text{crit}} - T_0}{\tau_i},
\]

where \( f_{\text{drag}}^* = Z_i^2 e^4 n_L / (4\pi e_n^2 m_e v_i^2) \). If we further make the assumption that \( T_{\text{crit}} \gg T_0 \) then this can be re-expressed as,

\[
\Xi_h > 1,
\]

with

\[
\Xi_h = \frac{I_i \tau_i}{3e_n} \frac{Z_i^2 e^4 L}{v_i^3} \pi \epsilon_0 m_e m_i.
\]

where \( I_i \) is the energy flux (intensity) of the ion pulse, and \( m_i \) is the ion mass. Note that we relate the energy flux to the ion beam density via,

\[
n_i = \frac{2I_i}{m_i v_i^3}.
\]

If we fix \( Z_i, n_e, \) and \( m_i \) then we can plot \( \Xi_h \) as a function of \( I_i \tau_i \) and \( \epsilon_i \). This is done in figure 1 for protons and \( n_e = 10^{39} \text{ m}^{-3} \).

From figure 1, we can see that for this target density and ion, we really require \( I_i \tau_i > 10^9 \text{ J m}^{-2} \) for proton energies around 1 MeV. This is equal to an energy flux of \( 10^{16} \text{ W cm}^{-2} \) for 10 ps.

We can thus examine the prospects for laser-generated proton beams. If we assume a laser to proton conversion efficiency of 5%, then this is achieved for \( I_i = 2 \times 10^{18} \text{ W cm}^{-2} \) and \( \tau_i = 1 \text{ ps} \). We can therefore see that relativistic laser intensities in the picosecond regime are well suited to entering the regime where a laser-generated proton beam can self-heat the target to the point of suppressing the drag. This of course neglects any spreading out the proton beam, as should be expected for a target that is distant from the primary (irradiated) target. However it may be a more reasonable estimate in the case where the proton beam rapidly enters the target of interest, such as is the case for front-surface acceleration of ions. It should also be noted that the above estimate does not just imply that relativistic short-pulses are able to access this regime, as intensities of \( 10^{16} \text{ W cm}^{-2} \) could likewise do so given \( \tau_i > 100 \text{ ps} \) (and a conversion efficiency of 10%). We can therefore conclude that the possibility of self-heating to the point of drag suppression is a serious possibility for modern laser systems.

Although we have obtained \( \Xi_h \) via estimation, we can make a more formal statement about it in the limit where we assume a ‘rigid beam’ (i.e. \( v_i \) and \( n_i \) are constants and the temporal profile of the pulse is fixed). In this case, the temperature evolves due to equation (5) alone, and it can be transformed into a dimensionless equation by expressing time in units of \( \tau_i \) and temperature in units of \( T_{\text{crit}} \) (i.e. \( t = \tau_i \), and \( T_e = T_{\text{crit}} T_e \)). Once this is done, the transformed heating equation is,

\[
\frac{\partial T_e}{\partial t} = \frac{I_i \tau_i}{v_i^3} \frac{Z_i^2 e^4 L}{2 \pi e_n^2 m_e^2 m_i} G(v_i/v_{e,th}),
\]

so it can see that \( \Xi_h \) is indeed a dimensionless number that determines the behaviour of the system, at least in this limit.

5. Electric field associated with heating

If the plasma is strongly heated by the incident ion beam, then there will be an electric field associated with this, as from Ohm’s law we have, \( E = -\nabla P_e / e_n \). If we assume that the electron density is only weakly perturbed then we can take the electric field to be \( E = -\nabla T_e \). If now work in a 1D approximation, then we have,

\[
\frac{\partial T_e}{\partial t} = \frac{2 f_{\text{drag}} v_i n_i}{3e_n}.
\]

If we assume a steadily propagating beam (at \( v_i \)) then the electric field will be,

\[
|E| = \frac{2 f_{\text{drag}} n_i}{3e_n},
\]

with the estimate for the maximum field we can obtain being,

\[
|E_{\text{max}}| \approx \frac{Z_i^2 e^3 I_i L}{2 \pi e_n^2 m_e m_i v_i^2}.
\]

If we denote the force on an ion due to the drag as \( f_{\text{drag}}^i \), and the force due to this electric field as \( f_{\text{el}}^i \) then we see that we quickly obtain,
In the special case where the ion beam is mixed and consists of protons and another (dominant) species, then we have,

\[ \frac{f^p_{\text{drag}}}{f^i_{\text{drag}}} = \frac{2}{3} Z^2 \left( \frac{n_i}{n_e} \right), \]

and thus in either case we see that any mitigation of the drag depends on the ion charge of the dominant beam component, and the ratio of the density of the dominant beam species to the background electron density. In the case of protons the drag can only possibly be significantly mitigated in the case where the proton beam density approaches a large fraction of the background electron density. On the other hand, a moderate Z ion beam could experience significant mitigation of the ion drag for \( n_i/n_e \ll 0.2 \). In the case of laser-generated ion beams interacting with a solid target, we generally expect the density in the ion beam to be significantly less than the target electron density (although there are certain regimes where high ion beam density can occur, e.g. [25, 26]). Therefore we do not expect this electric field effect to be significant in the majority of cases.

6. Electric fields associated with ion drag

There is another contribution to the electric field that needs to be considered. As there is a net drag force acting on the ion beam, there must be an equal but opposite force acting on the background electron fluid. This body force must be equal to \( f_{\text{drag}} \hat{b} \), where \( \hat{b} \) is a unit vector parallel with the direction of the ion beam. This additional term in the generalized Ohm’s law has, of course, been recognized before—see, for example [27]. For the case where the electron fluid is at rest, this implies that the Ohm’s law can be written as,

\[ E = -\nabla P_e + \frac{f_{\text{drag}}}{e} \frac{n_i}{n_e} \hat{b}, \]

in the absence of any pressure gradients, we still retain the contribution which is unique to ion drag, i.e.

\[ E = \frac{f_{\text{drag}}}{e} \frac{n_i}{n_e} \hat{b}, \]

which is close to, but somewhat larger than that obtained for the temperature gradient in the preceding section. Both effects can act simultaneously and reinforce one another. In terms of the importance of these field in mitigating ion drag, the same conclusion is reached: significant mitigation is only possible for ion beam densities that reach a substantial fraction of the target electron density. For laser-generated ion beams and solid density targets, this condition is unlikely to be achieved in the majority of cases, although it may be possible in certain alternative scenarios.

7. Implications for magnetic field generation

A prerequisite for magnetic field generation in the 2D/3D problem is the existence of sufficiently strong electric field. Although the electric fields discussed in sections 5 and 6 are not sufficient to mitigate the drag under those circumstances that we have argued will be typical, this does not mean that they could be negligible as far as magnetic field generation is concerned. From the induction equation, \( \partial_t B = -\nabla \times E \), we note that the magnitude of the flux density can be estimated to be,

\[ |B| \approx \frac{|E| \tau_i}{R_i}, \]

where \( R_i \) is the radius of the ion beam. If we use equation (12) from section 5 and we take \( I_1 = 10^{21} \text{ W m}^{-2}, v_1 = 10^7 \text{ m s}^{-1} \), \( Z = 1 \), and \( n_i = n_p \) then we arrive at \( E_{\text{max}} \approx 2 \times 10^8 \text{ V m}^{-1} \). If we then take \( \tau_i = 10 \text{ ps} \), and \( R_i = 10 \mu \text{m} \), then we would conclude that we should expect to see magnetic fields reaching flux densities around 100–200 T.

8. Numerical calculations

The previous discussion of the self-heating and electric field effects makes a number of approximations and simplifications. Naturally, it is interesting to see what happens if we relax these. To do this, we have employed a 1D hybrid code (APSIM). This treats the high energy ion beam kinetically, the target ions as a fluid, and the electrons are treated as a neutralizing massless fluid. This has certain similarities to a code previously developed by Sherlock et al [28], certain types of ‘hybrid PIC’ codes [29, 30], and is related to a simulation paradigm commonly used for ‘space plasmas’ [31]. The kinetic treatment of the ion beam is resolved by a macroparticle numerical model, as is the case in a traditional PIC code.

The kinetic ion macroparticles are thus evolved via the equations of motion:

\[ \frac{dv_{i,k}}{dt} = E_i, \]

and,

\[ \frac{dr_{i,k}}{dt} = Z_i e E + f_{\text{drag},i} \hat{v}_{i,k} + f_{\text{drag},q} \hat{v}_{i,k}, \]

where the ‘e’ and ‘i’ subscripts in the drag term indicate that both drag due to the fluid electrons and the fluid ions is included. In addition the kinetic ions are subject to angular scattering from the fluid ions.

The fluid ions evolve via hydrodynamic equations, i.e.

\[ \frac{\partial \rho}{\partial t} + \frac{\partial (\rho v_x)}{\partial x} = 0, \]

and,

\[ \frac{\partial v_x}{\partial t} + v_x \frac{\partial v_x}{\partial x} = -\frac{1}{\rho} \frac{\partial P_i}{\partial x} + Z_i e E, \]
as well as a fluid ion energy equation. The electrons are treated as a fluid that moves to maintain quasi-neutrality, i.e.

$$n_e = Z n_i^f + Z n_i^k,$$  \hspace{1cm} (24)

where $n_i^f$ and $n_i^k$ are the ion fluid density and kinetic ion density respectively, along with an electron energy equation to evolve the electron temperature. The electric field is determined by the Ohm’s law,

$$E = \frac{1}{\epsilon n_e} \nabla P_e.$$  \hspace{1cm} (25)

Note that we have opted to only include the thermobaric term, as the analysis in sections 5 and 6 indicated that we did not expect the electric fields to be sufficient to mitigate the drag, and so omitting a term of comparable magnitude should not affect the conclusions.

Open boundary conditions are used for the kinetic ions, and for the fluid ions we use open boundary conditions at one end (where kinetic ions are injected) and reflective boundary conditions at the other.

9. Results and discussion

Using aPSIS, we have carried out a sets of simulations to examine the issues raised in previous sections. Unless stated otherwise, we have used a grid with 1500 cells, and a cell size of $\Delta x = 0.1 \mu m$, giving a domain size of 150 $\mu m$. The time step was set to $\Delta t = 2$ fs. A total of $1.2 \times 10^9$ macroparticles were used for the fast ions, with 120 macroparticles being injected per time step. In the first set (set I), we considered a proton beam incident on a pure boron plasma at a uniform ion density of $n_i = 10^{29}$ m$^{-3}$, with $Z^* = 5$, and an initial electron and ion temperature of 100 eV. The proton beam is a monoenergetic beam with $v_i = 10^7$ m s$^{-1}$ (520 keV). The runs in the set are labelled A–E and we have tabulated pertinent input parameters in table 1, along with the nominal value of $\eta_h$ for each simulation. The simulations in set I were all run up to 30 ps.

In figure 2 the electron temperature profile is plotted at 90 ps for runs A–E. These plots clearly indicate that the ion beam penetrates to a much greater depth as $\eta_h$ increases above 1. This can be confirmed by determining another metric in each simulation, $R_{1/2}$, which we define to be the distance at which the last (macro)particle to fall to half its initial kinetic energy does so. This is plotted for runs A–E against the ion energy flux in figure 3, and against $\eta_h$ in figure 4.

Figures 3 and 4 confirm that the penetration range of ions increases with the ion beam energy flux and $\eta_h$. This is only possible if either the ion beam can heat the target rapidly enough to reduce the electron drag or if the electric field associated with the heating front can sufficiently mitigate the drag. In order to disentangle these two effects, a second set of simulations were carried out. In this set, set II, the same simulation parameters were employed, except that the effect of the electric field on the kinetic ions was turned off (along with the accompanying correction to the electron fluid energy). We will refer to the runs in this set as ‘II-x’ (with x denoting the specific run A–E). It was found that the changes to the simulation results were quite minimal.

In figure 5 the electron temperature profiles for simulations II-A and II-B are compared to their set I counterparts, and it can be readily seen that the differences are small. This is quite expected for this simulation based on the predictions made in section 5, as $n_i = 2I_i/(m_i v_i^2)$, so $n_i/n_e \sim 0.002$ in run A and II-A, and thus the force due to the electric field associated with the plasma heating should be small in this case. This reinforces
Figure 4. $R_{1/2}$ in runs A–E, in terms of $\Xi_h$. Note that range of a 500 keV proton in cold amorphous carbon at 2 g cc$^{-1}$ is $\approx 5 \mu$m (PSTAR NIST [32]).

Figure 5. Heating profiles at 30 ps in runs II-A and II-B as compared to runs A and B (see legends).

The conclusions reached at the end of sections 5 and 6, namely that the electric fields are unlikely to be important as long as we are in a regime where $n_i/n_e \ll 1$. As previously discussed, this regime is the more likely one for most laser-driven ion acceleration scenarios, although the $n_i \sim n_e$ scenario may be possible in more speculative scenarios (e.g. radiation pressure acceleration (RPA) based cases [25, 33]). Note that this largely justifies the omission of the drag term from the Ohm’s law employed in the code.

Although the electric field is not sufficient to mitigate the drag in these simulations this does not mean that electric fields are absent or weak. In figure 6 the electric field profiles in simulation I-A at 10 ps, 20 ps, and 30 ps are shown. It can be seen that the peak electric field does indeed reach around $10^8$ V m$^{-1}$, which is in accordance with the prediction of equation (12), and the discussion in sections 5 and 7.

However, in terms of the implications for magnetic field generation, the electric field profile indicates that, when only the thermobaric term is included in Ohms law, there is only a narrow spike in the electric field with a corresponding duration of about 1 ps, which means that the expected magnetic field generation would be considerably less than that anticipated in
This conclusion could change considerably with the inclusion of other terms in Ohm’s law, but this is a matter for future studies.

The next aspect that was investigated was the dependence of the temperature profile on the ion beam velocity. Since the drag heating starts to cut off when \( v_e \approx v_t \), this implies that higher temperatures will be reached for higher ion beam velocities. The counter-argument to this is that \( \Xi_h \propto v_i^{-6} \). To test this we carried out simulation set III (see table 2). Run III-A is identical to run I-A, and runs III-B and III-C are also identical except that \( v_i \) was set to \( 1.41 \times 10^7 \) m s\(^{-1} \) and \( 1.75 \times 10^7 \) m s\(^{-1} \) respectively. The temperature profiles at 30 ps are shown in figure 7 below.

As can be seen from figure 7, higher temperatures are reaches deeper into the target, although lower temperatures are achieved below 70 \( \mu \)m. Note that the values of \( \Xi_h \) in III-B and III-C are much lower, being 11.2 and 3.1 respectively. In run III-D we used \( v_i = 1.41 \times 10^7 \) m s\(^{-1} \) and \( I_i = 2.82 \times 10^{21} \) W m\(^{-2} \) (\( \Xi_h = 31.8 \)). In run III-E we used \( v_i = 1.75 \times 10^7 \) m s\(^{-1} \) and \( I_i = 5.2 \times 10^{21} \) W m\(^{-2} \) (\( \Xi_h = 15.9 \)). In these runs, \( \Xi_h \) is not so dramatically reduced compared to run III-A. The temperature profiles at 30 ps are plotted in figure 8, which shows that once \( \Xi_h \) is not so strongly reduced, a much hotter temperature is achieved throughout the target. If the interest in ion beam transport is more oriented towards heating a target (rather than being concerned with ion range for nuclear reactions), then this point could be pertinent.

The effect of angular scattering of the proton beam from the background ions was examined in simulation set IV, which was a repeat of simulation set I but with angular scattering removed. Although the difference in terms of the ion beam distribution are quite apparent, the temperature profile and the penetration range (\( R_{1/2} \)) was only modestly different, and we thus found that angular scattering was not strongly limiting the ion beam propagation, at least in this case.

As the increase in ion penetration that we have observed is due to target heating, there are obvious questions about how various cooling processes would affect these results. The most obvious are radiative cooling and thermal conduction. Radiative cooling via bremsstrahlung, assuming that all radiated power is lost, can be estimated via standard formulae, e.g.

\[
P_{\text{brems}} \approx 1.7 \times 10^{-40} Z^2 \sqrt{T_n n_e} \text{W m}^{-2},
\]

where \( T \) is in K. This gives an estimated cooling time of over 150 ps, which is much longer than the time-scales we have considered. Thermal conduction could be a concern in the 2D/3D problem, which, along with other 2D/3D issues, goes beyond the scope of this paper.

Finally, we acknowledge that there are studies which indicate that ‘warm’ targets can exhibit enhanced stopping of fast

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**Table 2.** Simulation parameters used in set III.

| Run   | Ion energy flux (W m\(^{-2} \)) | \( v_i \) (m s\(^{-1} \)) | \( \Xi_h \) |
|-------|---------------------------------|-------------------|-------------|
| III-A | \( 10^{21} \)                  | \( 10^7 \)        | 88.4        |
| III-B | \( 10^{21} \)                  | \( 1.41 \times 10^7 \) | 11.2        |
| III-C | \( 10^{21} \)                  | \( 1.75 \times 10^7 \) | 3.1         |
| III-D | \( 2.82 \times 10^{21} \)      | \( 1.41 \times 10^7 \) | 31.8        |
| III-E | \( 5.2 \times 10^{21} \)       | \( 1.75 \times 10^7 \) | 15.9        |
charged particles, e.g. the study of Zylstra et al [35]. We do not regard this as contradicting the results presented in this paper, as in Zylstra’s study the ‘warm’ plasma in question had a temperature in the 20–40 eV range. This is far too cool for drag reduction due to elevated electron temperature. The drag enhancement in this case is attributed to ionization potential depression, which makes it easier for a fast charged particle to excite bound electrons. We thus regard the findings of Zylstra as being of no relevance to those presented in this paper, as the two are considered to be completely different regimes.

10. Conclusions

In this paper, we have put forward the case that the combined ion–electron drag that is present when an intense (laser-generated) ion beam propagates into a target can become so extreme that the penetration range of the ions is greatly increased compared to the single ion range in the same target. We have derived an analytic metric for the onset of this regime in 1D hybrid ion-kinetic simulations. The regime can be entered for proton beam energy fluxes around $10^{20}$ W m$^{-2}$ ($10^{16}$ W cm$^{-2}$) with $\sim 10$ ps pulse durations. With reasonable assumptions about laser-to-proton energy conversion, it would therefore appear that laser-generated proton beams should be able to access this regime. For laser-generated ion beams that are produced by the standard target normal sheath acceleration (TNSA) mechanism, it is unlikely that the electric field will be sufficient to have a significant effect on the longitudinal dynamics of the ions which are dominated by the ion–electron drag.

These results show that it is possible for laser-generated proton and ion beams to ‘defy’ what one expects were one to treat as having only the range a single ion would have in a cold target, and we suggest that this possibility should be considered by fellow researchers in future. As we discussed in the introduction, there are a number of studies, which appear to require some mechanism that can greatly enhance the range of energetic ions in order to explain experimentally observed reaction yields. The enhancement in range required is typically 10 or more, but does vary from case to case. In the results presented here, we have shown in numerical simulations that the particle range can increase by a factor of 11 relative to the cold range. The enhancement of range due to the drag heating induced by the ion beam itself could therefore account for some of the observed yields in these experiments. In a similar fashion, the possibility of enhanced ranges may mean that laser-generated proton and ion beams have unique potential in applications involving beam-driven nuclear reactions, although future studies will be needed to determine if this is truly the case.

Data availability statement

The data that support the findings of this study are available upon reasonable request from the authors.

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ORCID ID

A P L Robinson https://orcid.org/0000-0002-3967-7647

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