Enhanced ion acceleration in the ultra-intense laser driven magnetized collisionless shocks

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Abstract
The effect of an externally applied magnetic field on the ion acceleration by laser-driven collisionless shocks is examined by means of multi-dimensional particle-in-cell simulations. For the interaction of ultra-intense sub-picosecond laser pulses with the near-relativistic critical-density plasma, the longitudinal transport of the laser generated fast electrons are significantly inhibited by the kilo-Tesla (kT) level transverse magnetic field, resulting in a thermal pressure which significantly exceeds the laser radiation pressure in the hot electron accumulation region. As a result, the accumulated plasma expands into the vacuum and leads to acceleration of a supersonic plasma flow in the opposite direction through the rocket effect, which streams into the target and drives a supercritical magnetized collisionless shock. In comparison with the case without the external magnetic field, where an electrostatic collisionless shock can be driven, the energy flux of the shock accelerated quasi-monoenergetic ion beam is considerably increased by an order of magnitude due to the strength enhancement of the magnetized shock.

1. Introduction
Laser-driven ion acceleration is one of the most active research fields in the laser-plasma physics due to its advantage in accelerating ions to very high (multi-MeV) energies over sub-millimeter distances [1, 2]. To date, several ion acceleration mechanisms have been proposed and observed experimentally, which include the target normal sheath acceleration [3–7], break-out afterburner [8–11], radiation pressure acceleration [12–18], collisionless shock acceleration (CSA) [19–29], etc. Recently, the CSA mechanism has attracted much attention due to its potential in generating monoenergetic ion beams in the µm-scale plasma targets, which is important for their applications in both science and medicine, such as the cancer therapy [30], proton radiography [31], and fast ignition inertial-confinement-fusion (ICF) [32].

In the most commonly studied CSA mechanism, the ultra-intense laser pulse first acts as a ‘piston’ to drive a supersonic unmagnetized plasma flow into the target through the hole-boring effect [33]. Then a high-Mach number electrostatic collisionless shock (also termed as the ion-acoustic shock) can be generated when the nonlinear steepening of the ion acoustic wave is balanced by the dissipation and dispersion effects as the flow streaming forward [34]. With the propagation of the shock, a portion of the upstream ions can be reflected back and accelerated to twice the shock velocity by the charge separation field at the shock front. Therefore, monoenergetic ion beams can be accelerated when the velocity of the electrostatic shock is kept steady by carefully tuning the laser-plasma parameters [24, 25].

However, despite the above merits, the CSA mechanism has suffered disadvantages in the energy transfer efficiency from laser to the shock accelerated ions, as shown in the recent experiments [26]. Recently, with the development of kT-level quasi-static magnetic fields in the laboratory [35], Mima et al have proposed a method
to improve the energy transfer efficiency by altering the laser-driven shock properties via externally applied longitudinal magnetic fields [27]. They have found that due to the anomalous deposition of the laser-generated fast electrons, electromagnetic fields which can randomly scatter the upstream ions can be generated at the shock front. As a result, the energy transfer efficiency is increased by a factor of \( \sim 6 \) due to the increase of the accelerated ion number compared with the case without the external magnetic field.

The ion acceleration efficiency in the CSA mechanism is strongly dependent on the shock strength, thus finding an effective way to enhance the shock strength is very important. In this paper, the effect of an externally applied transverse magnetic field on the laser-driven collisionless shock ion acceleration is studied via two-dimensional (2D) particle-in-cell (PIC) simulations. It is found that due to the external magnetic field, the laser-generated fast electrons are confined near the laser-plasma interaction (LPI) region, which results in a thermal pressure that significantly exceeds the laser radiation pressure. The fast expansion of the accumulated hot electrons then leads to the acceleration of a supersonic magnetized plasma flow into the target. As a consequence, a powerful magnetized collisionless shock is created with the shock strength considerably exceeding the electrostatic shock, which is driven without the external magnetic field. Compared with the electrostatic shock, both the energy and density of the shock accelerated ion beam is considerably enhanced in the magnetized shock.

2. Possibility for enhanced collisionless shock generation with an external magnetic field

Ion acceleration through reflection at the shock front occurs in collisionless shocks when the plasma cannot be fully dissipated by the shock. Here, the dissipation is defined as the process through which the kinetic energy of the upstream ion flow (in the shock frame) is dissipated into heat. In supercritical collisionless shocks, a part of the upstream ions at the tail of the distribution function that cannot overcome the potential barrier at the shock front (i.e. cannot be dissipated into heat) are reflected back and accelerated into the upstream region [36]. Thus the ion acceleration efficiency is increased with the enhancement of the shock strength due to the enhanced potential barrier [37]. In laser-driven collisionless shocks, the shock strength (e.g. shock velocity and Mach number) is dependent on the piston (hole-boring [33]) velocity, which is related to the plasma density \( n_p \) via \( v_{\text{piston}} \sim (1/n_p)^{1/2} \). On the other hand, the plasma density should be large enough to stay opaque to the incident laser pulse. Therefore, the near-relativistic-critical-density (NRCD) plasma is preferred for the laser-driven shock ion acceleration. However, it has been shown that the collisionless shock ion acceleration efficiency remains low even in the NRCD plasma [27].

In this paper, we consider the possibility of improving the ion acceleration efficiency via a kT-level transverse magnetic field, which is expected to alter the shock property by changing the fast electron transport dynamics. As schematically illustrated in figure 1, in the case without \( B_{0,z} \), the fast electrons freely stream into the target to heat up the plasma volume, and the shock velocity is determined by the light radiation pressure \( P_{\text{rad}} \) through the hole-boring effect. While in the case with \( B_{0,z} \), the fast electron transport is significantly inhibited resulting in a thermal pressure in the electron accumulation region of \( P_e \approx n_e k_B T_e \), where \( n_e \) and \( T_e \) are the density and temperature of the accumulated fast electrons, respectively. Since \( P_e \gg P_{\text{rad}} \) can be achieved taking into account of the electron accumulation and re-acceleration if proper magnetic field strength is chosen, a faster plasma flow can be generated via the fast expansion of the accumulated plasma compared with that through the hole-boring effect. In addition, the ion acoustic speed in the upstream region remains low due to the inhibition of the fast electron transport. Therefore, an enhanced collisionless shock with a larger velocity and larger Mach number is

Figure 1. Schematic presentation of the transport of the laser-generated fast electrons in the plasma target with (a) and without (b) an external transverse (perpendicular to the laser propagation direction) magnetic field \( B_{0,z} \). Collisionless shocks are expected to be driven in both cases but with significantly different properties due to the change of fast electron dynamics, as discussed in the text.
expected to be generated. Note that in this regime the NRCD plasma is also preferred due to the high electron acceleration efficiency and the low plasma inertia, which is favored for the powerful shock creation. In the next section, we will check the feasibility of this scheme by multi-dimensional PIC simulations in order to take into account the numerous nonlinear and kinetic effects self-consistently.

3. 2D PIC simulations of the laser-driven collisionless shock ion acceleration

We use the 2D3V PIC code ASCENT [38] to study the interaction of the ultra-intense laser pulse with the NRCD carbon plasma. The simulation box is $50 \lambda_L \times 8 \lambda_L$ large with (5000, 8000) uniform grids in the $(x, y)$ direction, where $\lambda_L = 1.06 \, \mu m$ is the laser wavelength. A uniform carbon ($C^{6+}$) plasma with an initial electron density of $n_{e0} = 8n_i$ exists in the simulation box with the initial vacuum-plasma interface at $x = 5 \lambda_L$, where $n_i \equiv \omega_i^2 m_i / 4\pi e^2$ is the critical electron plasma density and $\omega_i$ is the laser frequency. The initial temperatures of the electrons and ions are 2.5 keV and 0.25 keV, respectively. A $p$-polarized laser pulse with a uniform intensity profile in the $y$-direction enters the simulation box from the left boundary with a sharp rising edge, which keeps a constant intensity with $I_0 = 5 \times 10^{19}$ W cm$^{-2}$ after $5T_L$, where $T_L = \lambda_L / c$ is the laser period. In the simulations, 36 particles of each species are set per cell. For both fields and particles, we use absorbing boundary conditions in the $x$ direction and periodic boundary conditions in the $y$ direction. In addition, a fourth-order interpolation scheme is used to evaluate fields and currents.

We first present the PIC simulation results for two cases: the magnetized case with an external magnetic field $B_{0z} = 5.0$ kT and the non-magnetized case with $B_{0z} = 0$ kT. We note that the ions in the case with $B_{0z} = 5.0$ kT can be regarded as unmagnetized due to that the simulation time is less than $\Omega_i^{-1} \sim 4$ ps, where $\Omega_i$ is the ion gyro-frequency. In figure 2, snapshots of the particle phase-space distributions at $t = 200 T_L$ are present. Stable collisionless shock waves can be observed for both cases at the late time shown when a significant ion acceleration occurs, as exhibited in figures 2(b) and (d). In the non-magnetized case, a typical electrostatic shock structure develops when the laser-driven plasma flow streams into the target. Here, the role of the laser pulse is two-fold: it generates fast electrons through the $J \times B$ mechanism at a frequency of $2\Omega_i$ [39] to heat up the plasma volume (figure 2(a)), and acts as a piston to drive a supersonic plasma flow with the velocity of $v_{\text{piston}} = \left(1 + \eta f \right) I_0 / mn_i c^{3/2} \sim 0.037c$ [33], where $\eta f \approx 0.4$ is the laser reflection efficiency measured from the simulation, $m_i$ and $n_i$ are the carbon ion mass and density, respectively (here, the notations in the non-magnetized case are supplied by the symbol $'$. That is, $v_{\text{piston}}$ is larger than the ion acoustic speed.
accumulated hot electrons. As illustrated in Figure 2, the laser-generated fast electrons are constrained by the external magnetic field instead of collisions due to their long mean free path. While in the magnetized case, the transport of the laser-accelerated fast electrons are significantly inhibited by the external magnetic field, as shown in Figure 2(c). These fast electrons have a typical gyro-radius of [41]:

\[ r_e = \left( \frac{m_e c^2}{e B_{0,z}} \right) (\gamma_e^2 - 1)^{1/2} = 1.7 \times 10^3 (\gamma_e^2 - 1)^{1/2} B_{0,z}^{-1} \text{ cm}, \]

where \( \gamma_e \) is the relativistic factor and \( B_{0,z} \) is in gauss (1 T = 10^4 gauss). Note that since \( B_{0,z} \) is much less than the magnetic field amplitude of the laser, the \( J \times B \) mechanism is still the dominant electron acceleration mechanism. Therefore, \( \gamma_e \) can be estimated by the ponderomotive scaling law [33]: \( \gamma_e = (1 + \alpha_l^2/2)^{1/2} \). Then the electron gyro-radius can be estimated to be \( r_e = 1.7 \times 10^3 a_l B_{0,z}^{-1} / 2 \approx 1 \mu\text{m} \). Therefore, most of the laser-generated fast electrons are confined near the LPI region instead of propagating forward, resulting in a large thermal pressure in the electron accumulation region. For example, the density and effective temperature of the hot electrons in 18 < x/\( \lambda_L \) < 19 at t = 200 \( T_L \) are \( \approx 10 n_e \) and \( \approx 5.5 \text{ MeV} \), respectively, where the effective electron temperature is determined by fitting the simulated electron energy spectrum with the Maxwellian distribution. Then the electron thermal pressure can be estimated to be \( P_e \approx n_e k_B T_e \approx 100 \text{ Gbar} \), which is much larger than the laser radiation pressure \( P_{\text{rad}} = (1 + \eta) I_L / c \approx 20 \text{ Gbar} \), where \( \eta \approx 0.2 \) is the laser reflection efficiency. The reduced laser reflection efficiency (or enhanced absorption efficiency) in the magnetized case can be attributed to the more pronounced density perturbation and the re-acceleration of the recirculating fast electrons, which also results in a higher hot electron temperature. Note that although the recirculating electrons can be accelerated to a couple times the ponderomotive potential, they can still be effectively confined due to the compression of the magnetic field with the accumulation of the plasma (figure 4(d)).

Since \( P_e \gg P_{\text{rad}} \), a fast plasma flow is accelerated into the plasma target with the fast expansion of the accumulated hot electrons. As illustrated in figure 3(b), due to the momentum conservation, the plasma flow is generated through the rocket effect when the accumulated plasma blows out into the vacuum. The physical mechanism is similar to that of the non-relativistic laser-driven ablation [42], except that the laser heated electrons are stopped by the external magnetic field instead of collisions due to their long mean free path. Consequently, the laser absorption region is not separated from the ablation front by an electron heat conduction region but sited near the ablation front with a positive velocity. Another important difference is that in the non-relativistic regime the light pressure cannot cause significant hydrodynamic motion of the plasma.
While in the relativistic regime, the light pressure is large enough to drive a collisionless shock wave through the hole-boring effect. Therefore, in order for the ablation (or blow-out) driven mechanism to be dominant, the condition of $P_i \gg P_{\text{rad}}$ should be satisfied. In addition, the blow-out process shows a periodicity with the period of $\sim 12 T_i$, as shown in figure 3(b), which is about two times of the typical electron gyro-period ($2\pi \sigma_e/c \approx 6 T_i$), indicating that a coherent acceleration, accumulation and expansion process of the fast electrons exists.

The plasma flow velocity driven by the rocket effect is much larger than that driven by the hole-boring effect in the non-magnetized case, as shown in figures 3(a) and (b), which also significantly exceeds the magnetosonic speed $c_{\text{ms}} = (\sqrt{\gamma_A + \gamma_i})^{1/2} \approx 0.03c$, where $v_{\Lambda} = B_0 z/(4\pi n_i m_i)^{1/2}$ is the Alfvén speed and $c_i = (Z_K B_0 T_e/m_i)^{1/2}$ is the ion acoustic speed. Note that the electron temperature in the far upstream region of the magnetized shock remains nearly unchanged ($T_{ei} = 2.5$ keV) due to the inhibition of the fast electrons. As the supersonic magnetized plasma flow stream into the target, a high-Mach magnetized collisionless shock wave (which exhibits a laminar structure) can be driven through the wave steepening process that is balanced by the dissipation effect [43]. Note that the magnetized shock is created at the time scale of an order of the electron gyro-period instead of the ion gyro-period which is the typical formation time of the piston driven magnetized shock [44]. This is mainly due to that a large potential drop is generated between the downstream and upstream regions at the very beginning of the shock formation, which is already strong enough to accelerate the upstream ions and thus can significantly boost the shock formation.

In the non-magnetized case, the electrostatic collisionless shock has a shock velocity of $v_{sh}' \approx 0.048c$ (determined by the hole-boring velocity) with the Mach number of $M' = v_{sh}'/c_i' \approx 1.6$, which is larger than the critical Mach number of the electrostatic shock ($\sim 1.6$). While in the magnetized case, the magnetized collisionless shock propagates at the velocity of $v_{sh} \approx 0.08c$ (determined by the ablation velocity) with the magnetosonic Mach number of $M = v_{sh}/c_{\text{ms}} \approx 26.7$, which far exceeds the critical Mach number of the magnetized shock ($\sim 3$) [36]. Here, a quantitative prediction of the ablation velocity is difficult due to the complex LPI and electron transport dynamics. For these supercritical collisionless shocks, whether electrostatic or magnetized, ion acceleration can occur as an additional energy dissipation since the plasma cannot be fully dissipated by the shock. This is illustrated in figures 3(a) and (b), where typical trajectories of the shock accelerated ions are shown by the imposed blue dashed lines. Two main features can be identified: (i) the upstream ions are accelerated via reflection at the shock front to about twice of the shock velocity, and (ii) the ion acceleration efficiency in the magnetized shock is much larger than that of the electrostatic shock.

In both cases, the shock propagates at nearly a constant velocity after its formation. Therefore, quasi-monoenergetic ion beams can be accelerated, as shown in figure 3(d). In the non-magnetized case, the shock accelerated ions have a low energy spread of $\sim 20\%$ (FWHM) with the ion energy centered at $\sim 47$ MeV. While in the magnetized case, the energy spread of the shock accelerated ions remains low but the averaged ion energy is significantly increased to $\sim 141$ MeV. That is, the energy of the shock accelerated ions is increased by a factor of $\sim 3$ due to the enhancement of the shock strength in the magnetized case. The strength enhancement of the magnetized shock also results in an increase of the ion beam density by a factor of $\sim 6$. Therefore, the energy flux of the shock accelerated ions is considerably increased by an order of magnitude in the magnetized case. Here, the ion acceleration efficiency is quantified as the ratio of the energy flux of the shock accelerated ion beam to the laser pulse, which is $\sim 0.6\%$ in the non-magnetized case and $\sim 10\%$ in the magnetized case.

A direct result of the strength enhancement of the magnetized shock is the enhanced charge separation electric field at the shock front. As shown in figure 4(c), the longitudinal electric field localized at the magnetized shock front can reach a very large amplitude of $\sim 20$ TV m$^{-1}$, which is about 4 times larger than that of the electrostatic shock (figure 4(a)). Correspondingly, a larger electrostatic potential drop is generated at the magnetized shock front. As shown in figure 3(c), the potential difference between the downstream and the far upstream region in the magnetized shock is $e \Delta \Phi \approx 6$ MeV, which is much larger than that of the non-magnetized case $e \Delta \Phi' \approx 2$ MeV. This is the main reason for the increase of the energy flux of the shock accelerated ions in the magnetized shock. In addition, a train of trailing ion-acoustic waves is generated in the downstream region of the electrostatic shock due to the dispersion effect [34], as illustrated in figures 3(c) and 4(a), which lead to the generation of ion phase-space holes [45] due to ion trapping (figure 2(b)) when entering the nonlinear stage. However, the ion trapping is not significant in the magnetized shock due to the deformed downstream region.

In addition to the electric field, the magnetic field $B_z$ distribution also exhibits very different features in these two cases, as illustrated in figure 4. In the non-magnetized case, filament magnetic field structures can be observed in both the downstream and upstream regions due to the cumulative effect of the electron thermal anisotropy and the counter-streaming electron beams [46-48]. While in the magnetized case, the compression of the magnetic field in the shock downstream region is the main feature. The magnetic field is compressed to $\sim 3 B_0 z$ in the downstream region of the magnetized shock (figure 4(d)), indicating that the shock is very powerful with a large compression ratio. In addition, note that the gyro-radius of the shock accelerated ions in the upstream region of the magnetized shock is $r_i = v_{th}/\omega_{ci} \sim 200$ $\mu$m, where $v_{th} \sim 0.16c$ is the ion velocity and...
ωci = ZeB0, z/m_e c is the ion gyro-frequency. Therefore, the longitudinal transport of these ions is not significantly affected by the external magnetic field since r_i is much larger than the distance traveled by the ions.

Having clarified the enhancement of the laser-driven shock ion acceleration by the externally applied magnetic field, we proceed to discuss the influence of magnetic field strength on the collisionless shock property and the subsequent ion acceleration process. Another two cases are studied with B_z,0 = 2.0 kT and 10.0 kT, while other simulation parameters are kept unchanged. The energy density distributions of electrons for both cases are shown in figure 5. For B_z,0 = 2.0 kT, the laser-generated fast electrons cannot be effectively confined by the external magnetic field due to their large gyro-radius (figure 5(a)), which results in a large plasma beta (\beta_e \gg 1, \beta_e is the ratio of the electron plasma pressure to the magnetic pressure) in the upstream region. As a result, the shock property is similar to the non-magnetized case, and the energy of the shock accelerated ions is only slightly increased in comparison with the electrostatic shock (figure 3(d)). While for B_z,0 = 10.0 kT, the shock property and the energy of the shock accelerated ions is similar to the case with B_z,0 = 5.0 kT, except for a further increase of the accelerated ion beam density (figures 3(d) and 5(d)). However, considering the difficulty of generating larger quasi-static magnetic field in the laboratory, as well as the potential

Figure 4. The electric field E_x and magnetic field B_z distributions at t = 200 T_i for (a), (b) B_z,0 = 0 kT and (c), (d) 5.0 kT. The unit of the electromagnetic (EM) fields is \textit{m}_e \textit{ω}_L / e (1 unit = 3 TV m\(^{-1}\) for electric field and 10.0 kT for magnetic field). The black dashed line indicates the position of shock front.

Figure 5. The energy density distribution of hot electrons with the energy between 0.5 and 5.0 MeV (normalized by m_e c^2 n_i) at t = 200 T_i for (a) B_z,0 = 2.0 kT and (b) B_z,0 = 10.0 kT. The white dashed line indicates the position of shock front.
side effect of the larger magnetic field on the ion transport, $B_{0,z} = 5.0$ kT can be considered as a more suitable choice of the magnetic field for the unique set of laser-plasma parameters we studied.

In order to provide an insight into the three-dimensional aspect of this problem, an s-polarized simulation is carried out with the polarization vector out of the simulation $(x,y)$ plane and an external magnetic field with $B_{0,y} = 5.0$ kT. The s-polarized simulation exhibits a qualitatively similar feature with the $p$-polarized simulation, except for a reduced electron heating efficiency (figure 6(a)) [49, 50], which results in a reduced shock strength and more complex shock structure. The complicated shock structure is mainly due to the Rayleigh–Taylor (RT) instability developed near the laser-plasma interface, as illustrated in figure 7(a), which is driven when the blow-out (or ablated) low-density plasma accelerates the dense plasma inward [42]. In contrast, the shock wave maintains a quasi-one dimensional structure in the $p$-polarized case (figure 7(b)) due to the enhanced ablative stabilization effect at the higher ablation velocity [51]. Consequently, the quality of the shock accelerated ion beam is somewhat deteriorated in the $s$-polarized simulation, as illustrated in figure 6(b), including an increase of the ion energy spread to $\sim 35\%$ (FWHM) and a decrease of the ion energy flux by $\sim 50\%$ compared with the $p$-polarized simulation (with $B_{0,z} = 5.0$ kT). Nevertheless, the ion acceleration efficiency in the $s$-polarized simulation is still considerably enhanced in comparison with the electrostatic shock.

Finally, we discuss the requirement on the temporal-spatial profiles of the laser pulse for the efficient magnetized shock creation and ion acceleration. For the temporal profile, a sub-picosecond laser pulse is preferred considering the detrimental effects on the shock velocity at short pulse duration due to the fast time-varying laser intensity, and on the shock structure at long duration due to the possible development of RT instabilities. While for the spatial profile, a large laser spot size is preferred with $d_{z} \gg r_{e} \approx 1 \mu$m for an efficient confinement of the fast electrons taking into account of the transverse leaking of fast electrons from the LPI region, despite that the leaking effect can be compensated by the rapidly newborn fast electrons (e.g. the $J \times B$ mechanism generates fast electrons at a frequency of $2\omega_{pe}$). Last but not least, we want to comment on the possibility of driving the powerful collisionless shock through the blow-out regime with the assistance of the self-
generated quasi-static magnetic field in the LPI, which has been found to be strong enough to confine the fast electrons near the plasma surface [52].

4. Summary and discussion

To conclude, we have proposed a novel method to enhance the laser-driven collisionless shock ion acceleration via an externally applied transverse magnetic field. It is found through 2D PIC simulations that the external magnetic field results in a distinct change in both the property and strength of the laser-driven collisionless shock. An enhanced magnetized shock is driven through the blow-out effect of the accumulated fast electrons due to the external magnetic field, instead of the electrostatic shock driven through the hole-boring effect when the external magnetic field is absent. As a consequence, both the energy and density of the shock accelerated quasi-monoenergetic ion beam are considerably enhanced, leading to an increase of the ion energy flux by an order of magnitude. In addition to the application in laser-driven ion acceleration, generating strong magnetized collisionless shocks in the laboratory with this method can also promote our understanding of some important space and astrophysical phenomenons, such as the cosmic ray acceleration [53].

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References

[1] Daido H, Nishibuchi M and Pirozhkov A S 2012 Rep. Prog. Phys. 75 056401
[2] Macchi A, Borghesi M and Passoni M 2013 Rev. Mod. Phys. 85 731
[3] Snively R et al 2000 Phys. Rev. Lett. 85 2945
[4] Wilks S C, Langdon A B, Cowan T E, Roth M, Singh M, Hatchett S, Key M H, Pennington D, MacKinnon A and Snively R A 2001 Phys. Plasmas 8 542–9
[5] Esirkepov T Z et al 2002 Phys. Rev. Lett. 89 175003
[6] Fuchs J et al 2006 Nat. Phys. 2 48
[7] Arefiev A, Toncian T and Fiksel G 2016 New J. Phys. 18 105011
[8] Yin L, Albright B, Hegelich B and Fernández Á 2006 Laser Part. Beams 24 291–8
[9] Henig A et al 2009 Phys. Rev. Lett. 103 045002
[10] Hegelich B et al 2013 New J. Phys. 15 085015
[11] Jung D et al 2013 New J. Phys. 15 023007
[12] Esirkepov T, Borghesi M, Bulanov S, Mourou G and Tajima T 2004 Phys. Rev. Lett. 92 175003
[13] Robinson A, Zepf M, Kar S, Evans R and Bélier C 2008 New J. Phys. 10 013021
[14] Qiao B, Zepf M, Borghesi M and Geissler M 2009 Phys. Rev. Lett. 102 145002
[15] Henig A et al 2009 Phys. Rev. Lett. 103 245003
[16] Qiao B, Zepf M, Borghesi M, Dromey B, Geissler M, Karmakar A and Gibbon P 2010 Phys. Rev. Lett. 105 155002
[17] Yu T P, Pukhov A, Silvets G and Chen M 2010 Phys. Rev. Lett. 105 065002
[18] Palmer C A et al 2011 Phys. Rev. Lett. 106 014801
[19] Denavit J 1992 Phys. Rev. Lett. 69 3052
[20] Silva L O, Martí M, Davies J R, Fonseca R A, Ren C, Tsung F S and Mori W B 2004 Phys. Rev. Lett. 92 015002
[21] Chen M, Sheng Z M, Dong Q L, He M Q, Weng S M, Li Y T and Zhang J 2007 Phys. Plasma 14 113106
[22] Zhang X, Shen B, Yu M, Li X, Jin Z, Wang F and Wen M 2007 Phys. Plasma 14 113108
[23] Fiuza F, Fonseca R, Tonge J, Mori W and Silva L 2012 Phys. Rev. Lett. 108 235004
[24] Fiuza F, Stockem A, Boella E, Fonseca R A, Silva L O, Haberberger D, Tochitsky S, Gong C, Mori W B and Joshi C 2012 Phys. Rev. Lett. 109 215001
[25] Fiuza F, Stockem A, Boella E, Fonseca R, Silva L, Haberberger D, Tochitsky S, Mori W and Joshi C 2013 Phys. Plasmas 20 056304
[26] Haberberger D, Tochitsky S, Fiuza F, Gong C, Fonseca R A, Silva L O, Mori W B and Joshi C 2012 Nat. Phys. 8 95
[27] Mima K, Jia Q, Cai H B, Taguchi T, Nagatomo H, Sanz J R and Honrubia J 2016 J. Phys.: Conf. Ser. 717 012070
[28] Zhang W L, Qiao B, Shen X F, You W Y, Huang T W, Yan X Q, Wu S Z, Zhou C T and He X T 2016 New J. Phys. 18 093029
[29] Zhang H et al 2017 Phys. Rev. Lett. 119 166801
[30] Bulanov S, Esirkepov T, Khoroshkov V, Kuznetsov A and Pegoraro F 2002 Phys. Lett. A 299 240–7
[31] Borghesi M et al 2002 Phys. Plasmas 9 2214–20
[32] Roth M et al 2001 Phys. Rev. Lett. 86 436–9
[33] Wilks S C, Kruer W L, Tabak M and Langdon A B 1992 Phys. Rev. Lett. 69 1383–6
[34] Sagdeev R Z 1966 Rev. Plasma Phys. 4 23
[35] Santos J et al 2015 New J. Phys. 17 083051
[36] Balogh A and Treumann R A 2013 Physics of Collisionless Shock Waves (New York: Springer)
[37] Forslund D and Shonk C 1970 Phys. Rev. Lett. 25 1699
[38] Cai H B, Mima K, Zhou W M, Jozaki T, Nagatomo H, Sunahara A and Mason R J 2009 Phys. Rev. Lett. 102 245001
[39] Kruer W and Estabrook K 1985 Phys. Fluids 28 430–2
[40] May J, Tonge J, Fiuza F, Fonseca R A, Silva L O, Ren C and Mori W B 2011 Phys. Rev. E 84 025401
[41] Huia J 2004 NRL: Plasma Formulary (Washington, DC: Naval Research Laboratory)
[42] Atzeni S and Meyer-ter Vehn J 2004 The Physics of Inertial Fusion: Beam-Plasma Interaction, Hydrodynamics, Hot Dense Matter (New York: Oxford University Press)
[43] Tidman D A and Krall N A 1971 Shock Waves in Collisionless Plasmas (New York: Wiley)
[44] Schaeffer D, Fox W, Haberberger D, Fiksel G, Bhattacharjee A, Barnak D, Hu S and Garmanchewski K 2017 Phys. Rev. Lett. 119 025001
[45] Eliasson B and Shukla P K 2006 Phys. Rep. 422 225–90
[46] Bret A, Gremlillet L and Dieckmann M E 2010 Phys. Plasmas 17 120501
[47] Stockem A, Fiuza F, Bret A, Fonseca R and Silva L 2014 Sci. Rep. 4 3934
[48] Stockem A, Grasmayer T, Fonseca R and Silva L 2014 Phys. Rev. Lett. 113 105002
[49] Ren C, Tsoufras M, Tsung F, Mori W, Amorini S, Fonseca R, Silva L, Adam J and Heron A 2004 Phys. Rev. Lett. 93 185004
[50] Stark D J, Yin L, Albright B J and Guo F 2017 Phys. Plasmas 24 053103
[51] Takabe H, Mima K, Montierth L and Morse R 1985 Phys. Fluids 28 3676–82
[52] Bigongiari A, Raynaud M, Riconda C, Héron A and Macchi A 2011 Phys. Plasmas 18 102701
[53] Reville B, Bell A and Gregori G 2013 New J. Phys. 15 015015