The controllable electron-heating by external magnetic fields at relativistic laser-solid interactions in the presence of large scale pre-plasmas

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Abstract

The two-stage electron acceleration/heating model (Wu et al 2017 Nucl. Fusion 57 016007 and Wu et al 2016 Phys. Plasmas 23 123116) is extended to the study of laser magnetized-plasmas interactions at relativistic intensities and in the presence of large-scale preformed plasmas. It is shown that the electron-heating efficiency is a controllable value by the external magnetic fields. Detailed studies indicate that for a right-hand circularly polarized laser, the electron heating efficiency depends on both strength and directions of external magnetic fields. The electron-heating is dramatically enhanced when the external magnetic field is of $B < 0$. When magnetic field is of negative direction, i.e. $B < 0$, it trends to suppress the electron heating. The underlying physics—the dependences of electron-heating on both the strength and directions of the external magnetic fields—is uncovered. With $-\infty < B < 1$, the electron-heating is explained by the synergetic effects by longitudinal charge separation electric field and the reflected ‘envelop-modulated’ CP laser. It is indicated that the ‘modulation depth’ of reflected CP laser is significantly determined by the external magnetic fields, which will in turn influence the efficiency of the electron-heating. While with $B > 1$, a laser front sharpening mechanism is identified at relativistic laser magnetized-plasmas interactions, which is responsible for the dramatical enhancement of electron-heating.

Keywords: laser plasma interactions, laser driven ion acceleration, fast ignition of inertial confinement fusion

(Some figures may appear in colour only in the online journal)

1. Introduction

Laser-matter interactions at relativistic intensities have attracted great attention from both experimental and theoretical investigations, because of its significant effects on a number of applications, such as laser driven ion acceleration [1–7], and fast ignition of inertial confinement fusion [8–14]. Although the contrast ratio of the start-of-art laser beams can be as high as $10^{-8}$ [10], considerable pre-plasma can still build up in front of a solid density target. The pre-formed
plasma always exists in laser-matter interaction at relativistic intensities, therefore, understanding electron-heating at intense laser-solid interactions in the present of large scale pre-plasmas [15–20] is important to a number of great applications.

In a recent work, a two-stage electron-acceleration (heating) model was proposed to understand intense laser-solid interactions influenced by large scale pre-plasmas [22, 23]. The electron-heating is attributed to the synergetic effects by the longitudinal charge separation electric field $E_z$ and the ponderomotive force of the conflicting laser beams. Efficiency of the electron-heating depends on both the pre-plasma scale-length and the laser intensity. Maximal possible energy gain within this synergetic process was analyzed, with value of $\Delta \varepsilon \sim (IL_p)^{3/2}$, where $I$ is laser intensity and $L_p$ is pre-plasma scale-length.

Recently, extremely strong magnetic field up to 1500 T and with ns duration is obtained by using a novel capacitor-coil target design making use of the hot electrons generated during the laser-target interactions [24]. The unprecedented strong and long-lasting magnetic fields have aroused great interests in the community [25–33]. The properties of laser-plasmas interactions will be quite different in the presence of external magnetic fields, which can be applied to enhance the energy coupling in laser plasma interactions. For example, the laser propagation in highly magnetized high-density plasmas has been investigated. It is shown that due to the whistler regime [31], moderately intense circularly polarized (CP) laser light can enter and propagate in high-density plasma and heat it efficiently. Furthermore, the control the transmission of laser light in over-dense plasma has been achieved by applied magnetic field pulses [32], which is able to produce intense only two-cycle light pulses as well as long-lasting self-magnetic fields. Except external magnetic fields, the self-generated magnetic fields, which usually produced in intense laser solid interactions, could also strongly influence the electron dynamics. For example, when a fast electron beam propagates through a solid target with resistivity gradient, resistive magnetic field could occur and inhibit the beam propagation. This effect can have consequences on the laser-driven ion acceleration [34] and Bremsstrahlung photon emission [35].

A new fast ignition concept—magnetically assisted-fast ignition—was proposed recently [28], which is also based on the unprecedented strong and long-lasting magnetic fields. For magnetically assisted fast ignition, one need to figure out how do external magnetic fields influence (i) the hot electron generation, (ii) the proration of hot electrons in solid target and (iii) the energy deposition of hot electrons. In addition, the magnetically assisted-laser driven ion acceleration concept [29], was also proposed recently, with the resulting proton beam much better collimated compared to a beam generated without applying a KT-level magnetic field. It has been well recognized that electron acceleration in the laser plasma interaction at the front surface of the target is the key to effective proton acceleration. In this paper, we have investigated laser magnetized-plasmas interactions at relativistic intensities and in the presence of large-scale preformed plasmas. The present work, therefore, is strongly related to the proposed magnetically assisted-fast ignition [28] and -laser driven ion accelerations concept [29].

It is shown that the electron-heating efficiency can be controlled by both the strength and directions of the external magnetic fields, which is applied along the laser propagation direction. Detailed studies indicate that for a right-hand (RH) CP laser, the electron heating efficiency depends on both strength and directions of external magnetic fields. The electron-heating is dramatically enhanced when the external magnetic field is of $B \equiv \omega_e / \omega_0 > 1$, and magnetic field of negative direction, i.e. $B < 0$, trends to suppress the electron heating, where $\omega_e$ and $\omega_0$ are cyclotron frequency of electrons in magnetic fields and laser frequency in vacuum. The underlining physics—the dependences of electron-heating on both the strength and directions of the external magnetic fields—is uncovered. Within $-\infty < B < 1$, the variation of electron-heating is explained by the synergetic acceleration by longitudinal charge separation electric field and the reflected ‘envelop-modulated’ CP laser. It is indicated that the ‘modulation depth’ is significantly determined by the external magnetic fields, which will in turn influence the efficiency of the synergetic acceleration. While with $B > 1$, a laser front sharpening mechanism is identified at relativistic laser magnetized-plasmas interactions, which is responsible for the dramatical enhancement of electron-heating.

This paper is arranged as follows: one-dimensional simulations are performed in section 2. Theoretical analysis of the simulation results are demonstrated in section 3. In section 4, two-dimensional (2D) simulations are performed to further confirm our findings. The conclusions and discussions are given in section 5.

2. One-dimensional simulations

One-dimensional particle-in-cell (PIC) code, LAPINE [36], are performed to model laser magnetized-plasmas interactions, where weighted-particle and 4th order particle-cloud (FDTD) methods are applied to reduce the numerical noises. The schematic of simulation set-up is shown in figure 1. The size of the simulation box is 400 $\mu$m, with the region $0 < z < 100 \mu$m left as vacuum, and initial plasma density profile is taken as $n_e = n_{\text{solid}}/(1 + \exp[-2(z - z_0)/L_p])$, where $n_{\text{solid}} = 50n_{0}$, the solid plasma density, $z_0 = 180 \mu$m and $L_p = 10 \mu$m is the pre-plasma scale-length. The simulation box is divided into 40,000 grids, with each grid containing 100 electrons and ions. As we have done previously, to avoid the influences of ion mobility, the movement of ions is deliberately turn off. The simulation time step is 0.01 fs. The laser of intensity $10^{20}$ W cm$^{-2}$ with wavelength of 1 \mu m enters the simulation box from the left boundary, and along the laser propagation direction, an uniform magnetic field is applied. As shown in figure 1, we have placed two diagnostic planes to temporally record the electrons passing through. The first one is located at $z = 100 \mu$m recording electrons passing through at $z$-direction, and the other one is at $z = 300 \mu$m recording electrons of $z$-direction passed.
3. Theoretical explanations

From simulation results, we have found that electron energy spectra heavily depend on the strength and directions of the external magnetic fields, especially a dramatical enhancement of electron heating is observed when \( B > 1 \) is satisfied. In figure 3, we have presented the space–time evolution and phase-space plots of electrons under three typical external magnetic fields, with \( B = -5 \), \( B = 0 \) and \( B = 5 \). By the comparison with figures 3(a)–(c), we can find that the case with \( B = 5 \) is significantly different from the other two cases. From figure 3(c), we can identify the source of these super-high energetic electrons, which come from the very first electron bunch initially accelerated backward by the charge separation electric field. Furthermore, the very first electron bunch as shown in figures 3(e1) and (e2) seems to be of highly collimated.

In the following theoretical analysis, we will firstly study the gradual increase of electron heating efficiency with the increase of magnetic fields, \(-\infty < B < 1\), and then the special case with \( B > 1 \).

With \(-\infty < B < 1\) cases—When the incident laser arrives at the critical density surface and is reflected back, due to the formation of the steep interface of electron density, a strong delta-like charge separation field or the step-like electrostatic potential [21] is build up therein, which is strong enough to drive electrons to very high velocity within very short time and short length. This ‘initial large velocity’ can significantly simplify our following analysis. Imagine we are standing on the frame of a backward propagating electron, we will find that the incident laser pulse is oscillating very fast, and its only contribution to the motions of the electron is to increase its mass by a factor \( \gamma = (1 + \alpha^2)^{1/2} \) in an average way, however the reflected laser pulse is oscillating so slow that this electron can be captured and continually be accelerated backward by its ponderomotive force. If reflected laser pulse is of CP, the synergetic effects [22, 23] by the longitudinal charge separation electric field \( E_z \) and the ponderomotive force of the reflected laser beam is of no effect at all. However, in a recent work [23], we have found that the envelop of the reflected CP laser is significantly modulated. Once the ponderomotive force of modulated CP laser is non-zero, the synergetic acceleration is turned on again. The modulation of the reflected CP laser is due to the oscillation of the reflection point at critical-densities. When electrons leave the reflection point, charge separation field therein is lowered and when returning back, it is increased again. In reality, the oscillation period of reflection point is attributed to the circulating time of electrons. The dominant oscillation period is determined by the circulating time of the majority of electrons [23], i.e., \( \omega_{pe} \) of the emitting electrons. The density of the emitting electrons can be calculated by an energy flux balance equation. Therefore, the oscillation frequency is estimated to be \( \omega_0 / \sqrt{\gamma} \), where \( \omega_0 \) is the incident laser frequency in vacuum.

From figure 4, we can see that, in the presence of external magnetic fields, (a) \( B = -5 \), (b) \( B = -3 \) and (c) \( B = 0.5 \), the reflected CP laser is also modulated. Although the modulation
The modulated CP laser is typically of the following form, \( a_\text{e} \sim \sin(k_0(z - ct))[1 + \alpha \cos(k_m(z - ct))] \) and \( a_\text{r} \sim \cos(k_0(z - ct))[1 + \alpha \cos(k_m(z - ct))] \), with the ‘modulation depth’ defined as \( \alpha \equiv (a_{\text{max}} - a_{\text{min}})/(a_{\text{max}} + a_{\text{min}}) \). Recalling that it is the oscillation of the reflected laser beam, by the Doppler effect, the frequency on the mirror’s frame of reference becomes \( \omega = \omega_0[1 + v_{\text{ref}}(t)] \). When reflected back by the mirror, the reflected frequency becomes \( \omega = \omega_0[1 + v_{\text{ref}}(t)]^2 \sim \omega_0 + 2\omega_0v_{\text{ref}}(t) \). Here the velocity of the mirror is considered to be much smaller than the light speed, \( v_{\text{ref}}(t) \ll 1 \). As \( \alpha \) is a Lorentz invariant, for incident laser of \( E_\text{e} \sim \sin(\omega_0 t) \) and \( E_\text{r} \sim \cos(\omega_0 t) \), the reflected laser becomes, \( E_\text{e} \sim [1 + 2v_{\text{ref}}(t)] \sin(\omega_0 t) + 2v_{\text{ref}}(t)\omega_0 t] \sim [1 + 2v_{\text{ref}}(t)]\sin(\omega_0 t) + 2v_{\text{ref}}(t)\omega_0 t] \sim [1 + 2v_{\text{ref}}(t)]\cos(\omega_0 t) \). We immediately noticed that the ‘modulation depth’ of the reflected CP laser is determined by the oscillation velocity of the reflection point.

Actually the oscillation velocity of the reflection point heavily depends on the ponderomotive force acted on the electrons. Here we have, \( f_p \approx \frac{\hbar n_\text{e}(\text{bv})^2}{2} \), which is the energy balance equation of electrons initially located at the reflection point, where \( f_p \) is the ponderomotive force, \( \hbar \) is the distance departing from the origin, \( n_\text{e} \) is the number of electrons, and \( \text{bv} \) is the velocity of the oscillation. Taking into account \( T_\text{em} = \frac{\hbar n_\text{e}}{\text{bv}} \), we can immediately obtain \( \text{bv} \approx \sqrt{\frac{f_p}{n_\text{e}m_\text{e}}} \). Now we know, the ‘modulation depth’ of reflected CP laser depends on the oscillation velocity of reflection point, while the oscillation velocity depends on the ponderomotive force acted on the electrons. Below, we shall figure out how the ponderomotive force acted on the electrons is influenced by external magnetic fields.

Here we refer to Maxwell equations and the electrons’ equations of motion. The relevant Maxwell equations are

\[
\nabla \times E_L = -\frac{1}{c} \frac{\partial B_L}{\partial t},
\]

\[
\nabla \times B_L = \frac{4\pi}{c} J + \frac{1}{c} \frac{\partial E_L}{\partial t}.
\]

Assuming laser of the form \( \exp[i \mathbf{k} \mathbf{z} - i \omega t] \) and \( J = e_n \mathbf{p}_e / \gamma \), we can find

\[
\mathbf{p}_e = -\gamma \left( \frac{e_0}{\gamma^2} \right) E_L.
\]

The electron’s equations of motions are

\[
\frac{m_e}{c} \frac{\partial \mathbf{p}_e}{\partial t} = -e E_{\text{ex}} - e \mathbf{p}_e B_{\text{ex}} / \gamma,
\]

\[
\frac{m_e}{c} \frac{\partial \mathbf{p}_e}{\partial t} = -e E_{\text{ex}} + e \mathbf{p}_e B_{\text{ex}} / \gamma.
\]

Where \( B_{\text{ex}} \) is the external magnetic fields. When combine equation (3) with equation (4), and into account \( -i \mathbf{E}_{\text{ex}} \sim E_{\text{ex}} \), we have \( \mathbf{p}_e = (e/\hbar m_e) [1 + m_e \omega_c (\omega^2 - c^2 k^2)]/[4\pi e^2 n_e \omega_c] E_{\text{ex}}, \) where \( \omega_c = e B_{\text{ex}} / m_e \). Similarly, when combine equation (3) with equation (5), we have \( \mathbf{p}_r = (e/\hbar m_e) [1 + m_e \omega_c (\omega^2 - c^2 k^2)]/[4\pi e^2 n_e \omega_c] E_{\text{ex}} \).

The ponderomotive force, i.e., \( (e/\gamma) \mathbf{p}_e \times B_{\text{ex}} \), is expressed as,

\[
f_p \sim \frac{e^2}{i \gamma m_e} \left[ 1 + \left( \frac{\omega_c}{\omega} \right)^2 - \frac{c^2 k^2}{4\pi e^2 n_e / m_e} \right] (E_L \times B_{\text{ex}}).
\]
The ponderomotive force can be further simplified by implementing the dispersion relation of CP laser under magnetized plasmas,

$$k^2 = \frac{\omega^2}{c^2} \left( 1 - \frac{\omega^2}{\omega_e^2} \frac{1}{\gamma - \omega_i/\omega} \right)$$

When considering $E_L = (-m_e c/e)(\partial A/\partial t)$, the ponderomotive force can be re-organized as

$$f_p \sim \frac{m_e c^2}{2\gamma} \left( 1 + \frac{B}{\gamma - B} \right) \frac{\partial A^2}{\partial z},$$

where $B$ is defined as $B \equiv \omega_i/\omega$. From equation (8), we find that the ponderomotive force heavily depends on $B$. The coefficient, $1 + B/(\gamma - B)$, changing with $B$ is plotted in figures 4(d), one can see that it is increasing with the increase of $B$. From figures 4(a)–(c), the increase of ‘modulation depth’ with the increase external magnetic fields can be well understood by the increase of the coefficient, $1 + B/(\gamma - B)$, with $B$.

When the reflected CP laser is of envelop modulated, the synergetic acceleration is turned on again. Considering the electron propagates backward with high velocity, the only contribution of the incident wave is to increase the electron mass in an averaged way [22, 23]. As we have done recently [22, 23], start from the $z$-momentum and energy equation of electron,

$$\frac{d(\gamma v_z)}{dt} = -\frac{1}{2\gamma} \frac{\partial a_z^2}{\partial z} + E_z,$$

$$\frac{d\gamma}{dt} = \frac{1}{2\gamma} \frac{\partial a_z^2}{\partial t} + E_z v_z,$$

(9)

where $E_z$ is the charge separation field, the relativistic factor $\gamma$ is defined as $\gamma = \gamma_0 \gamma_1$, with $\gamma_1 = (1 + a^2/2 + a_c^2/\gamma^2)$, and $\gamma_0 = 1/(1 - v_c^2)^{1/2}$. The backward going wave $a_z$, based on figure 4(b), has the following form,

$$a_z \sim \sin(k_0(z - ct))(1 + \alpha \cos(k_m(z - ct))),$$

$$a_z \sim \cos(k_0(z - ct))(1 + \alpha \cos(k_m(z - ct))),$$

(10)

where $k_0$ is the wave number of incoming wave, $k_m$ is the wave number of the modulating wave, and $0 < \alpha < 1$ is the modulation depth, which is defined as $\alpha \equiv (a_{\max} - a_{\min})/(a_{\max} + a_{\min})$. Following equation (10) and defining $a^2 = Ra^2$ with $R$ of the reflection ratio, $\gamma_1$ can be written as $\gamma_1 = 1 + a^2/2 + \alpha a_c^2 + (1 - 2\alpha)$, $a^2 \sin^2(k_m/2)(z - ct)$, here $\alpha = (1/2)(1 - \alpha)/(1 + \alpha)$. From equations (9), we find

$$\frac{d}{dt}\gamma \gamma_1(1 - v_c) = -E_z(1 - v_c).$$

When assuming electric field $E_z$ is of constant, we can obtain $\gamma_1 \gamma_1(1 - v_c) = \sigma_{\gamma} E_z(1 - v_c)$, where $t_0$ is the time at which the electron crosses $z = z_0$ and $\sigma_{\gamma} = \gamma_1 \gamma_1(1 - v_c)$, $t = t_0$. Note for the highly relativistic case, we have $\sigma_{\gamma} \sim 1/(2) \gamma_1 \gamma_1 \gamma_1 \gamma_1 \ll 1$. The trajectory of the electron $z$ can be found by introducing a local time $\tau = k_0(t - z)$, with $dz/d\tau = v_c/(1 - v_c)$. Using $v_c$, $dz/d\tau$ can be found to be $dz/d\tau = 1/\gamma(z(z + \tau - z_0))$, where $f(\tau) = \gamma_1(\gamma_1 \gamma_1(1 - v_c))$. Making use of the inequality $\sigma_{\gamma} \ll 1, \sigma_{\gamma + \gamma_1} \ll 1$ and $E_z \tau \ll 1$, $\Delta \tau$ can then be rewritten as

$$\Delta \tau = \frac{1}{2} \int_0^{\tau} \frac{d\gamma}{\sigma_{\gamma} - E_z \tau} d\tau.$$

Taking into account, $k_m = k_0/2$, the maximal in-phase time is $\tau = 2\pi$, then the maximal-possible energy gain within limited length $L$ can be expressed as

$$\Delta E(2\pi) = \frac{a_e^2}{8E_z} (1 - 2\alpha)^2 f(\sigma_{\gamma}).$$

(12)

and

$$L = \frac{1}{2E_z} \left[ \gamma_1^2(2\pi + \tau) - \frac{\gamma_1^2(\tau)}{\sigma_{\gamma} - 2\pi} - \frac{a_e^2}{4f(\sigma_{\gamma})} \right] - \pi .$$

(13)
where \( \sigma_E = \sigma_b / E_t \geq 2\pi \), and

\[
f(\sigma_E) = \frac{1}{\sigma_E} \int_0^{2\pi} \frac{\sin(x/2)}{\cos(x/2)} dx.
\] (14)

Maximal-possible energy gain within limited longitudinal length \( L \) can be found, assuming \( a \gg 1 \) and \( L \gg 1 \),

\[
\Delta E = \frac{R}{8}(1 - 2a) f(\sigma_E) aL^{1/2} = \eta aL^{1/2},
\] (15)

with \( g^2(\sigma_E) = [(R - 2\alpha R)\sigma_E + \pi(2\alpha R + 1)] / [2\sigma_E(\sigma_E - 2\pi)] - R(\sigma_E) / 2 \). Coefficient \( \eta \) is a function of \( \sigma_E, R \) and \( \alpha \). \( L \) is determined by pre-plasma scale-length. As shown in figure 5, with the decrease of \( \alpha \equiv 1/2(1 - \alpha) / (1 + \alpha) \), \( L \) is increasing. This kind of behavior exactly explains the energy spectra shown in figure 2, with the increase of external magnetic field, the electron acceleration/heating is also increasing.

With \( B > 1 \) cases—As shown in figure 4(e), when \( B = 5 \), the amplitude of reflected CP laser is almost close to zero. Under such conditions, synergetic accelerations by charge separation field and ponderomotive force of reflected CP laser would never be the dominant contribution to the observed dramatic electron-heating.

According to equation (7), we find that for \( B \equiv \omega_b / \omega > 1 \), there is a resonant region with \( B < \gamma < B + \omega_b^2 / \omega^2 \). Within the resonant region, electromagnetic waves can not propagate. Instead, they are absorbed by plasmas. However for \( -\infty < B < 1 \), there is no stopping band for relativistic intense laser beams. Therefore, the physics with \( B > 1 \) is totally different with cases of \( -\infty < B < 1 \).

Figure 6 shows the \( v_b^2 - \gamma \) diagram following the dispersion relation in equation (7), where \( v_b^2 = \omega^2 / k^2c^2 \). From figure 6(a), we can see that for external magnetic field of \( B = 5 \), there is a resonant region with \( 5 < \gamma < 5.25 \). Beyond this resonant region and with the increase of \( \gamma \), we find \( v_b^2 \) is rapidly decreasing. For \( \gamma \gg 1 \), we have \( v_b \sim c^2 / \gamma_b \), which means \( v_b \) is rapidly increasing. Considering a Gaussian laser propagating through the magnetized plasmas, these two effects, ‘the resonant band’ and ‘group velocity speeding with the increase of \( \gamma \)’, would lead to the laser front sharpening. As the ponderomotive force is determined by \( \partial \gamma / \partial z \), the laser front sharpening effect would dramatically enhance the ponderomotive force of the laser beam. Compared with figure 3(c), it is clearly indicated that the laser front sharpening effect can be triggered only for \( \omega_b / \omega > 1 \) and with the direction of external magnetic field in \( z \)-direction.

To compensate the analysis above, we also run a serial of PIC simulations, where uniform background density of \( n_e = 0.2n_i \), and laser pulse of \( \sin^2 \gamma \) rising profile with rising period of \( 10T_0 \) are assumed. The laser is of RH-CP with intensity \( 10^{20} W \text{ cm}^{-2} \). We change the directions and strength of the external magnetic fields to study their influences on the evolution of the laser front. As shown in figure 7(a), for external magnetic field with \( B = 5 \), there are two effects undertaken for the laser front. Within the region having laser amplitude \( a \sim \gamma \sim 5 \), the electromagnetic waves are absorbed by plasmas due to the resonant motion of electrons. Beyond the absorbed region, we can also observe significant laser front congestion. These two effects, shown in figure 7(a), which collectively lead to the laser front sharpening, are well consistent with our theoretical analysis. In contrast, as shown in figure 7(b), if the external magnetic is of \( z \)-direction, we do not observe the laser front sharpening effect, as expected by the theoretical analysis.

From figure 7(a), at the resonant region \( a \sim \gamma \sim 5 \), where the laser is absorbed by plasmas, we find intense electron accumulation therein. These accumulated electrons should be the source of the highly energetic and highly collimated electron bunch, as shown in figures 3(c1) and (c2). The laser front sharpening effect would dramatically increase the laser ponderomotive force, because of \( f_\rho \sim \partial \gamma / \partial z \). This ponderomotive force would push all electrons forward,
leaving ions behind, until at a position where the charge separation induced electric force therein is equal to the ponderomotive force, i.e. $E_z \sim f_p$. As we know $E_z \sim n_e z$, we can easily conclude that the reflection position of the very first electron bunch depends on the external magnetic field, which is clearly confirmed by figure 3. For $B = 5$, the reflection position is at $z = 148 \mu m$, while for $B = -5$, it is around $z = 130 \mu m$.

The electron bunch will eventually be accelerated backward by the charge separation electric field. During this process, there exist a phase-locked mechanism which would result in the formation of highly collimated electron bunch. The profile of the charge separation field is shown with red line in figure 7(a). For electrons traveling faster, the electric field experienced by them is small, while for electrons traveling slower, the electric field acted on them is large. This kind of charge separation electric field profile will ‘phase-lock’ the electrons, which share the same physics with the phase stable acceleration of ions by laser radiation pressure [37], eventually leading to the formation of highly collimated electron bunch, as shown in figures 3(c1) and (c2).

The highly energetic and highly collimated electron bunch would build an intense electrostatic potential, as shown by the red curves in figure 2, with its peak energy several times larger than electron kinetic energy. Finally, electrons reflected by this potential barrier will acquire energy several times as large of their initial values.

4. 2D simulations

Here, 2D PIC simulations are performed to further confirm this mechanism. Here to avoid the extensive calculation burden, we use a smaller simulation box and shorter laser pulse duration. The simulation box is $40 \mu m \times 40 \mu m (L_x \times L_y)$, with grid size $\Delta z = 0.02 \mu m$ and $\Delta y = 0.1 \mu m$. The CP Gaussian laser is of intensity $10^{20} W cm^{-2}$, with laser wavelength of 1 $\mu m$, pulse duration of 100 fs and radius of 15 $\mu m$, with an external magnetic field of 15 000 T along $z$ directions, and with pre-plasma scale-length of 5 $\mu m$.

Figure 8 shows the field intensity and electron density distribution at $t = 50T_0$. Recalling it is the reflected modulated CP laser that initiate the acceleration, the ‘modulation depth’ of reflected CP laser heavily depends on the external magnetic fields. From the electron density distribution, figures 8(a1)–(c1), we can see that with the increase of magnetic fields, the backward acceleration of electrons is also increasing. Figure 9 shows the energy spectra of electrons collected by diagnostic plane located at $z = 50 \mu m$ (not shown) influenced by external magnetic fields, the heating efficiency is also increasing with the increase of external magnetic fields.

Note, the ‘cut-off’ energy produced in 2D simulations is much smaller than that in 1D. Recalling the scaling law of electron heating/acceleration efficiency is of $\sim AL^{1/2}$, this is partially due to the following reasons: (i) pre-plasmas scale-length used in 2D simulation is 5 $\mu m$ instead of 10 $\mu m$ as used in 1D simulation, therefore the cut-off energy should be smaller by $\sqrt{2}$ times; (ii) laser pulse used in 2D simulations is of Gaussian instead of flat-top as used in 1D simulations, therefore the ‘cut-off’ energy should also be reduced somewhat. Furthermore, transverse effects might also play roles. When $B > 1$, as we have analyzed, there exist a phase-locked mechanism which would result in the formation of highly collimated electron bunch. When compared with figure 3(c1) and figure 8(c2), electron bunch accelerated backward by the charge separation electric field is transversely disturbed. This kind of transverse perturbation would reduce the acceleration efficiency. As we know in the laser radiation pressure acceleration (RPA), there also exits transverse instabilities [38]. Ion acceleration efficiency produced in 2D/3D RPA simulation is also much smaller than 1D RPA.

5. Discussions and conclusions

In summary, the two-stage electron acceleration/heating model is extended to the study of laser magnetized-plasmas interactions at relativistic intensities and in the presence of large-scale preformed plasmas. It is shown that the electron-heating efficiency is a controllable value by the external magnetic fields. Detailed studies indicate that for a RH-CP laser, the electron heating efficiency is increasing with the increase of external magnetic fields. Especially, the electron-heating is dramatically enhanced when the external magnetic field is of $B \equiv \omega_c / \omega_0 > 1$. The underlying physics—the dependences of electron-heating on both the strength and directions of the external magnetic fields—is uncovered. With $-\infty < B < 1$, the electron-heating is explained by the synergetic effects by longitudinal charge separation electric field and the reflected ‘envelop-modulated’ CP laser. It is indicated that the ‘modulation depth’ is significantly determined by the external magnetic fields, which will in turn
influence the efficiency of the electron-heating. While with \( B > 1 \), a laser front sharpening mechanism is identified at relativistic laser magnetized-plasmas interactions, which is responsible for the dramatical enhancement of electron-heating. To further confirm the phenomena we observed, 2D PIC simulations are performed, with results consistent with the one from 1D simulations.

For Rd-glass laser with typical wavelength 1.0 \( \mu \text{m} \), the corresponding magnetic field can be as high as 10 000 T. Although the highest static magnetic field obtained in experiment is 1500 T, the laser front sharpening effect and the related phoneme we reported in this work can still be confirmed by experiment. We would recommend to use CO\(_2\) laser facility [39], because of its longer wavelength. Laser wavelength of CO\(_2\) laser is 10 \( \mu \text{m} \), which means the \( B \) can be as high as 1.5, considering the highest static magnetic field reported so far.

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**References**

[1] Bin J H et al 2015 Phys. Rev. Lett. 115 064801
[2] Wu D, Zheng C Y, Zhou C T, Yan X Q, Yu M Y and He X T 2013 Phys. Plasmas 20 023102
[3] Chen M, Pukhov A, Yu T P and Sheng Z M 2009 Phys. Rev. Lett. 103 024801
[4] Qiao B, Zept M, Borghesi M and Geissler M 2009 Phys. Rev. Lett. 102 145002
[5] Yan X Q, Wu H C, Sheng Z M, Chen J E and Meyer-ter Vehn J 2009 Phys. Rev. Lett. 103 135001
[6] Wu D, Zheng C Y, Qiao B, Zhou C T, Yan X Q, Yu M Y and He X T 2014 Phys. Rev. E 90 023101
[7] Wu D, Qiao B, McGuffey C, He X T and Beg F N 2014 Phys. Plasmas 21 123118
[8] Van Woerkom L et al 2008 Phys. Plasmas 15 056304
[9] Tabak M, Hammer J, Glinsky M E, Kruei W L, Wilks S C, Woodworth J, Campbell E M, Perry M D and Mason R J 1994 Phys. Plasmas 1 1626
[10] Ma T et al 2012 Phys. Rev. Lett. 108 115004
[11] MacPhee A G et al 2010 Phys. Rev. Lett. 104 055002
[12] Wu D, He X T, Yu W and Fritzsche S 2017 Phys. Rev. E 95 023207
[13] Wu D, He X T, Yu W and Fritzsche S 2017 Phys. Rev. E 95 023208
[14] He X T, Li J W, Fan Z F, Wang L F, Liu J, Lan K, Wu J F and Ye W H 2016 Phys. Plasmas 23 082706
[15] Paradkar B S, Wei M S, Yabuuchi T, Stephens R B, Haines M G, Krasheninnikov S I and Beg F N 2011 Phys. Rev. E 83 046401
[16] Paradkar B S, Krasheninnikov S I and Beg F N 2012 Phys. Plasmas 19 060703
[17] Krasheninnikov S I 2014 Phys. Plasmas 21 104510
[18] Sorokovikova A, Arefiyev A V, McGuffey C, Qiao B, Robinson A P L, Wei M S, McLean H S and Beg F N 2016 Phys. Rev. Lett. 116 155001
[19] Kemp A J, Sentoku Y and Tabak M 2009 Phys. Rev. E 79 066406
[20] Sheng Z M, Mima K, Zhang J and Meyer-ter-Vehn J 2004 Phys. Rev. E 69 016407
[21] Mishra R, Sentoku Y and Kemp A J 2009 Phys. Plasmas 16 1127
[22] Wu D, Krasheninnikov S I, Luan S X and Yu W 2017 Nucl. Fusion 57 016007
[23] Wu D, Krasheninnikov S I, Luan S X and Yu W 2016 Phys. Plasmas 23 123116
[24] Fujioka S et al 2013 Sci. Rep. 3 1170
[25] Zhu B J et al 2015 Appl. Phys. Lett. 107 261903
[26] Cai H B, Zhu S P and He X T 2013 Phys. Plasmas 20 072701
[27] Jia Q, Cai H B, Wang W W, Zhu S P, Sheng Z M and He X T 2013 Phys. Plasmas 20 032113
[28] Wang W M, Gibbon P, Sheng Z M and Li Y T 2015 Phys. Rev. Lett. 114 015001
[29] Arefiyev A, Toncian T and Fiksel G 2016 New J. Phys. 18 105011
[30] Yang X H, Yu W, Xu H, Yu M Y, Ge Z Y, Xu B B, Zhao H B, Ma Y Y, Shao F Q and Borghesi M 2015 Appl. Phys. Lett. 106 224103
[31] Luan S X, Yu W, Li F Y, Wu D, Sheng Z M, Yu M Y and Zhang J 2016 Phys. Rev. E 94 053207
[32] Ma G, Yu W, Yu M Y, Luan S X and Wu D 2016 Phys. Rev. E 93 053209
[33] Feng W, Li J Q and Kishimoto Y 2016 Phys. Plasmas 23 032102
[34] Gizzi L A, Betti S, Forster E, Giuliani D, Hofer S, Koster P, Labate L, Lotzsch R, Robinson A P L and Uschmann I 2011 Phys. Rev. ST Accel. Beams 14 011301
[35] Zamponi F et al 2010 Phys. Rev. Lett. 105 085001
[36] Xu H, Chang W W, Zhuo H B, Cao L H and Yue Z W 2002 Chin. J. Comput. Phys. 19 305
[37] Yan X Q, Lin C, Sheng Z M, Guo Z Y, Liu B C, Lu Y R, Fang J X and Chen J E 2008 Phys. Rev. Lett. 100 135003
[38] Yan X Q, Wu H C, Sheng Z M, Chen J E and Meyer-ter-Vehn J 2009 Phys. Rev. Lett. 103 135001
[39] Haberberger D, Tochitsky S, Fiuza F, Gong C, Fonseca R A, Silva L O, Mori W B and Joshi C 2011 Nat. Phys. 8 95

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