Erratum: Ion acceleration via ‘nonlinear vacuum heating’ by the laser pulse obliquely incident on a thin foil target (2016 Plasma Phys. Control. Fusion 58 025003)

A Yogo\textsuperscript{1,2}, S V Bulanov\textsuperscript{2}, M Mori\textsuperscript{2}, K Ogura\textsuperscript{2}, T Zh Esirkepov\textsuperscript{2}, A S Pirozhkov\textsuperscript{2}, M Kanasaki\textsuperscript{2}, H Sakaki\textsuperscript{2}, Y Fukuda\textsuperscript{2}, P R Bolton\textsuperscript{2}, H Nishimura\textsuperscript{1} and K Kondo\textsuperscript{2}

\textsuperscript{1} Institute of Laser Engineering, Osaka University, Suita, Osaka Prefecture 565-0871, Japan
\textsuperscript{2} Kansai Photon Science Institute, Japan Atomic Energy Agency, 8-1-7 Umemidai, Kizugawa, Kyoto 619-0215, Japan

E-mail: yogo-a@ile.osaka-u.ac.jp

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In the published version, equation (30) was incorrectly written as:

\[ \varepsilon_p \sim \frac{m_e c^2}{\cos^2 \phi} \left\{ \frac{2}{\pi} E \left[ A \sqrt{1 + A} \right] \right\} - 1 \]

The correct equation is as follows:

\[ \varepsilon_p \sim \frac{m_e c^2}{\cos^2 \phi} \left\{ \frac{2}{\pi} \sqrt{1 + A} E \left( \frac{A}{1 + A} \right) - 1 \right\} \]
Ion acceleration via ‘nonlinear vacuum heating’ by the laser pulse obliquely incident on a thin foil target

A Yogo, S V Bulanov, M Mori, K Ogura, T Zh Esirkepov, A S Pirozhkov, M Kanasaki, H Sakaki, Y Fukuda, P R Bolton, H Nishimura and K Kondo

1 Institute of Laser Engineering, Osaka University, Suita, Osaka Prefecture 565-0871, Japan
2 Kansai Photon Science Institute, Japan Atomic Energy Agency, 8-1-7 Umemidai, Kizugawa, Kyoto 619-0215, Japan

E-mail: yogo-a@ile.osaka-u.ac.jp

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Abstract

The dependence of the energy of ions accelerated during the interaction of the laser pulse obliquely incident on the thin foil target on the laser polarisation is studied experimentally and theoretically. We found that the ion energy being maximal for the p-polarisation drastically decreases when the pulse becomes s-polarised. The experimentally found dependencies of the ion energy are explained by invoking anomalous electron heating, which results in high electrostatic potential formation at the target surface. The anomalous heating of electrons beyond the energy of quiver motion in the laser field is described within the framework of a theoretical model of a driven oscillator with a step-like nonlinearity. We have demonstrated that the electron anomalous heating can be realised in two regimes: nonlinear resonance and stochastic heating, depending on the extent of stochasticity.

Keywords: high-intensity laser, laser ion acceleration, laser plasma interaction

1. Introduction

The discovery of laser-driven ion acceleration [1] has opened up a broad field of potential applications involving fast ignition [2, 3] and hadron therapy [4–6]. Most intensively investigated has been the mechanism of ion acceleration with charge separation (CS) fields generated on the thin foil target surfaces, including target normal sheath acceleration (TNSA) [1]. Nowadays, increasing attention is also paid to different mechanisms; however, the TNSA mechanism still remains crucial to understanding the physics of ion acceleration at moderate laser intensity regimes.

The generation of an ion accelerating CS field is governed by the absorption mechanism of laser energy into electrons. In the case of interactions between an obliquely incident laser and overdense plasma, one of the most predominant absorbed mechanism is the Brunel effect [7]. In this model, electrons on the boundary are directly accelerated by the electric field of the laser, when the electrons are accelerated to the normal direction of the surface with a momentum \( p_{\text{max}} = 2mc_0a_0 \) at maximum, where \( c \) is the speed of light in vacuum, \( m \) and \( e \) are the electron mass and charge, and \( a_0 = E_0/mc_0 \) is a dimensionless amplitude of the laser radiation, with \( E_0 \) and \( \omega \) being the laser electric field amplitude and frequency. It can be expressed in terms of the laser intensity \( I \) and wavelength \( \lambda = 2\pi c/\omega \) as \( a_0 = 0.85\sqrt{I/(10^{18} \text{ W cm}^{-2})}\lambda/(\mu\text{m}^2) \).

To explain the acceleration beyond \( p_{\text{max}} \), a nonlinear effect should be implemented into the analysis. D’yachenko and Imshennik [8] numerically revealed that electrons quivered by the laser field are kicked out by the potential boundary of the plasma and anomalously gain momentum depending on time (see also [11, 12]). Paradkar et al [9] and Krasheninnikov [10] theoretically showed stochastic heating of electrons by assuming a model dependent on the coordinate of the electric
field corresponding to the field in the vicinity of a thin foil target. The heating mechanism of electrons driven by \( j \times B \) force was precisely discussed by Mulser et al. [13], Bauer et al. [14] and Debayle et al. [15]. Using particle-in-cell simulations, Taguchi et al. [16] found that the heating of electrons transiting through a small size target can be higher than that by electron quivering, and called this phenomenon ‘nonlinear resonance absorption’. Robinson et al. [17] proposed that non-wake-field acceleration beyond the momentum limit is gas plasma.

In this manuscript, we develop the mechanism of ion acceleration involving the nonlinear effect on the electron energisation in order to describe the accelerated proton energy during the laser interaction with a thin foil target. By considering a model of a driven oscillator with a step-like nonlinearity, the growth of the chaotic behaviour of electrons is revealed. Here, we call the mechanism nonlinear vacuum heating. As a result, the experimental results on the proton energy and its dependency on the laser polarisation are successfully predicted.

2. Experiments

The experiment was performed using the JLITE-X laser at JAEGA, which delivers ultrashort (\( \tau = 40 \, \text{fs} \), FWHM) linearly polarised pulses with a central wavelength of \( \lambda = 800 \, \text{nm} \). The laser contrast was increased to 10^6 thanks to the Insertable Pulse Cleaning Module [19]. The beam is focused on a spot 8 \( \mu \text{m} \) (FWHM) in diameter using an \( f = 160 \, \text{mm} \) off-axis parabolic mirror (OAP) with an incidence angle of 45° on silicon-nitride foils. The polarisation angle is varied from \( p (\theta = 0^\circ) \) to \( s (\theta = 90^\circ) \) by using a half-waveplate located at the entrance of the OAP. The areal peak intensity obtained is \( I = 3.0 \times 10^{18} \, \text{W cm}^{-2} \), which corresponds to the dimensionless amplitude \( a_0 = 1.2 \). The pedestal intensity level in the sub-ns range is around \( 10^9 \, \text{W cm}^{-2} \), which is adequate to suppress the pre-plasma formation on the ultrathin foil surface. Accelerated protons are measured by a Time-of-Flight (TOF) online detector [18] located at the normal direction of the target rear surface. A typical energy spectrum obtained is shown in figure 1(b). It is confirmed using ion-track (CR-39) detection that the protons (\( >1.2 \, \text{MeV} \)) have a 15° angular divergence distributed along the normal of the target surface [19].

In the experiment we analyse the dependence of the energy of accelerated ions on the laser polarisation. Figure 1 shows the maximum energy of the protons (circles) observed for 500 nm thick targets as a function of the polarisation angle \( \theta \). Here, the proton energy drastically decreases as the polarisation is changed from \( p \) to \( s \) mode. A similar dependency was also observed [21] for a thicker (13 \( \mu \text{m} \)) target under high-contrast laser conditions. On the other hand, there is a clear difference from the case of the low-contrast laser when the proton energy for \( s \)-polarisation was only 30% lower [20] than that for \( p \)-polarisation. As we can see, the obvious dependency on the polarisation angle is characteristic of the interaction between steep-gradient plasma surfaces and obliquely-incident laser pulses. To understand the physics underlying the phenomenon, we discuss the motion of the electrons driven by a strong EM wave in the vicinity of the plasma-vacuum interface.

3. Oblique incidence of the laser radiation on a thin foil target

3.1. Electron quiver motion

Let us consider a plane electromagnetic (EM) wave interacting with the thin foil located in the plane \( x = 0 \). The laser angle of incidence equals \( \phi \) (see figure 2(a)). The EM wave is linearly polarised and its electric field \( \mathbf{E} \) has a polarisation angle \( \theta \) with respect to the \( x-y \) plane: \( \theta = 0^\circ \) for \( p \)-polarisation and \( \theta = 90^\circ \) for \( s \)-polarisation. In order to simplify the consideration, we perform a Lorentz transformation from the laboratory frame \( K \) to the frame of reference \( K' \) moving with velocity \( V = c \sin \phi \)
along the target surface (the y direction). In the K′ frame, the EM wave is incident normally onto the surface [24, 26], and the gamma factor associated with the velocity V is given by

\[ \gamma = \frac{1}{\sqrt{1 - \frac{V^2}{c^2}}} = \frac{1}{\cos \phi}. \]  

(1)

The Fourier component for a monochromatic plane wave in vacuum [22] is given in the K′ frame by equations

\[ \nabla \times E' = \frac{\omega'}{c} B', \]  

\[ \nabla \times B' = -\frac{\omega'}{c} E', \]  

(2)

(3)

where the wave frequency equals \( \omega' = \omega \gamma = \omega \cos \phi \). Using \( \nabla \cdot E' = 0 \), equations (2) and (3) can be written as follows

\[ \nabla^2 E' + \frac{\omega^2}{c^2} E' = 0. \]  

(4)

Using equations (2)–(4) we obtain the general solution of the Maxwell equations,

\[ E' = E_0 e^{i\omega' (x-x'/c) + \phi} + E_e e^{i\omega' (x+x'/c)}. \]  

(5)

with the electric field amplitudes of incident \( E_0 \) and reflected \( E_e \) waves in the K′ frame.

When the EM wave interacts with the plasma with a steep gradient satisfying the condition \( r_L / L \gg 1 \), we can use the Leontovich boundary conditions [23] for the tangential components of the electric and magnetic fields \( E_t \) and \( B_t \):

\[ E_t = \zeta B_t \times n, \]  

(6)

where \( n \) is the unit vector along the inward normal to the plasma surface and \( \zeta \) is the surface impedance operator. Here \( r_L \) is the electron quiver radius and \( L \) is the scale-length of the plasma inhomogeneity. Assuming that the plasma surface can be treated as an ideal conductor (\( \zeta = 0 \)), we obtain the boundary condition

\[ E'_{t=0} = E_t = 0. \]  

(7)

Then, the electric field is written as

\[ E' = E_0 [e^{i\omega' (x-x'/c)} + e^{i\omega' (x+x'/c)}]. \]  

(8)

The magnetic field is obtained using equations (2) and (8) as follows:

\[ B' = E_0 [e^{i\omega' (x-x'/c)} + e^{i\omega' (x+x'/c)}]. \]  

(9)

Therefore, the real part of the magnetic field at the surface \( x' = 0 \) is given by

\[ B_{x=0}' = (0, -2E_0 \sin \theta \cos \omega t', 2E_0 \sin \theta \cos \omega t'). \]  

(10)

In the boosted frame of reference, K′, the acceleration of electrons in the x′ direction is driven by the Lorentz force equal to \(-e/c)v \times B\), where the magnetic field is taken at the \( x = 0 \) plane, \( B = B_{x=0}' \), and the velocity of the electrons is equal to \( v = (0, -V, 0) \). As a result the equation of motion for electrons can be written as

\[ \frac{dp'}{dt'} = -e/c [v \times B_{x=0}'], \]  

(11)

yields for the x′-component of the electron momentum

\[ \frac{dp'}{dt'} = 2m_e V a t_0 \omega \sin \theta \cos \omega t'. \]  

(12)

Hence, the x′-component of the electron momentum depends on time as

\[ p_{x'} = 2m_e V a t_0 \omega \sin \theta \cos \omega t', \]  

(13)

where \( V = c \sin \phi \). We see that the electron momentum \( p_{x'} \) is below \( 2m_e c^2 t_0 \), which corresponds to the momentum limit for the electrons accelerated by the obliquely incident laser via the Brunel effect [7].

### 3.2 Nonlinear mechanism of anomalous electron heating

Within the framework of the nonlinear vacuum heating mechanism, the electrons quivered by the laser field are kicked out by the steep-like electric field at the plasma boundary. As a result, the EM field energy is irreversibly converted into the electron kinetic energy, thus enhancing the electron heating. To describe the nonlinear motion of electrons, we assume that a step-like CS electric field is formed near the ion layer with a positive electric charge.

In addition, we should consider that actual laser pulses always have some extent of pedestal. In our experiment, although the pedestal intensity is around \( 10^{10} \) W cm\(^{-2} \) at the time 500 ps before the main pulse, the pedestal at 10–100 ps before the main pulse ranges from \( 10^{10} \) to \( 10^{12} \) W cm\(^{-2} \) [19].
which exceeds the typical ionisation threshold of a solid target. Therefore, it is natural to consider that the density of the extremely thin foil target, 100–500 nm thick in our experiment, will be reduced by the pedestal just before the main pulse arrival. This phenomenon was also discussed in some of the literature [25] showing that the target is not broken completely, keeping a near-critical density. In this case, the target density is slightly higher than the critical density only at the thin region located at the centre of the foil. Thus, it is natural to assume that the electron fluid element is initially situated in the middle of the target.

Hence, the equations of electron motion can be written in the form

\[ \dot{p} + \varepsilon_p \text{sign}(x) = a_0 \cos \omega t, \quad (14) \]

\[ \dot{x} = \frac{p}{(1 + p^2)^{1/2}}, \quad (15) \]

Here the sign function is equal to \( \text{sign}(x) = 1 \) for \( x > 0 \) and \( \text{sign}(x) = -1 \) for \( x < 0 \), and a dot denotes a differentiation with respect to time, \( t \). The electron momentum is normalised on \( m_c \), the time is measured in units \( \omega^{-1} \) and the coordinate is normalised on \( c/\omega \).

The parameter

\[ \varepsilon_p = 2\pi n_{ef}^2 l/m_c \omega, \quad (16) \]

introduced in [26], is proportional to the charge separation electric field, \( E_0 = 2\pi n_{ef} l \) (see figure 2(1)), normalised on \( m_c \omega/e \), for a thin foil with an electric charge surface density equal to \( enol \).

The formulated model of a driven oscillator with a step-like nonlinearity (see also [27]) helps to determine and elucidate the conditions when the stochastic vacuum heating takes place. The three-dimensional model used bears the key features of the problem addressed in [9, 10, 28, 29], but in a significantly simpler mathematical approximation, because a full description of the charged particle interaction with electromagnetic waves implies a consideration of a seven-dimensional dynamical system.

While the electron trajectory crosses the \( x = 0 \) plane, the sign of the static electric field changes abruptly. This can be considered as a ‘collision’, during which the oscillating driver electric field produces the work changing the electron energy. Depending on the phase, when the ‘collision’ happens the work can be either positive or negative, resulting in the electron energy increasing or decreasing.

The enhancement of the energy transfer from the driver electric field to the electron is expected to occur when a resonant condition (for details see [27]) takes place:

\[ a_0 = \pi \varepsilon/2, \quad (17) \]

which can be equivalent to the relationship of \( E_0 \approx E_c \). The expression for the momentum maximum is found in [27] for two different nonlinear regimes: nonresonance and stochastic heating.

In a nonresonant resonance regime, the electrons gain a momentum \( \Delta p \approx a_0 \) during the passage of the plane \( x = 0 \). Then, the momentum grows proportionally to the square root of time, written in

\[ p_m \approx (a_0 t/2)^{1/2} = a_0(t/\pi)^{1/2} \quad (18) \]

for the regime of nonlinear resonance. In the case of stochastic heating, the momentum dependence on time has a character of diffusion with the diffusion coefficient

\[ D_{pp} = \frac{2\Delta p \Delta p}{\tau} = \frac{a_0^2 \varepsilon}{2p_m}, \quad (19) \]

This yields for the time dependence of the maximum electron momentum

\[ p_m \approx (a_0^2 t/2)^{1/3} = a_0(t/\pi)^{1/3}, \quad (20) \]

noticed in [9].

To analyse in more detail the properties of nonlinear oscillations we have integrated equations (14) and (15) numerically. For the regularisation of the singularity on the left-hand side of equation (14) we use instead the sign function \( \text{sign}(x) \) and the function \( \text{Tanh}(x/l) \) with the layer width \( l \) (see figure 2(c)) very small compared to the particle displacement: \( l \ll \Delta x \). Figures 3(a) and (b) show the result of the numerical calculation of the momentum and the coordinates of the electrons as functions of time obtained for the condition close to our experiment. We can see clearly that the oscillations are nonlinear with the momentum abruptly dropping and rising again. In the time duration of the laser pulse, 40 fs (\( t = 95 \) in figure 3), the momentum reaches \( 8 m_c \), which is a few times higher than the quiver momentum corresponding to the Brueell effect. The oscillation amplitude is limited both in the momentum and coordinate directions, as shown in figure 3(c). The trajectory tightly fills the domain in the phase plane demonstrating the ergodic property of the system under consideration. In figure 3(d) we plot a Poincaré section with the form of a finite thickness web, which is broadened due to the stochastic nature of the particle motion. The property of the particle to migrate on the phase plane along the web is typical for minimal chaos regimes [30]. In other words, chaotic behaviour has begun to emerge, but does not evolve completely under these conditions.

3.3. Ion acceleration

The nonlinear vacuum heating conjecture is equivalent to the assumption that the energy conversion from the EM field to electron energy depends on time according to the nonlinear mechanism introduced above. The average square of the \( x \)-component of the electron momentum \( p_{nvh,x} \) scales with a laser pulse duration proportional to the number of wave periods \( N \) for the two different regimes of nonlinear vacuum heating as:

\[ p_{nvh,x}^2 \propto \begin{cases} N p_{\pi}^2 & \text{nonlinear resonance} \cr N^{2/3} p_{\pi}^2 & \text{stochastic heating} \end{cases} \quad (21) \]
To estimate the typical energy gain of fast protons, we use the step-like CS field equal to $E_0 \tanh(x/l)$ in the interval

$$-l/2 - \Delta x \leq x \leq l/2 + \Delta x,$$

where $\Delta x$ is the electron displacement and $l$ is the thickness of the positively charged layer of the foil target (figure 2(c)). When the target thickness is thin enough to satisfy the condition $l \approx \Delta x$, the electric field $E_0$ can be expressed as

$$E_0 = 2\pi n_0 e l = 2\pi n_0 e \Delta x$$

(23)

with the ion charge density equal to $n_0$. Assuming that the protons can be treated as test particles initially located on the surface of the layer, the proton energy gain $E_p$ can be estimated to be of the order of the electron kinetic energy, $E_e$:

$$E_p \sim E_e \approx \frac{\Delta x}{2}$$

(24)

The above approximation is applicable when the movement of the electrons is not strongly relativistic, i.e. $\gamma \approx 1$. In the boosted frame of reference $K'$, the kinetic energy can be found to be given by

$$E_p' \sim E'_e = E'_\text{total} - m_e c^2$$

(25)

$$E'_\text{total} = \frac{\omega'}{2\pi} \int \frac{d\rho}{\sqrt{m^2 c^4 + p^2 c^2 + p_{\text{max}}^2 c^2}}$$

(26)

where

$$p_v = \gamma M m_e V = m_e \tan \phi.$$ 

(27)

From equations (13) and (21), it follows that

$$p_{\text{max},\phi} \sim N' p_v = 2N' m_e a_0 \sin \phi \cos \theta \sin \omega' t'$$

(28)

with the index $j$ equal to $j = 1/2$ for the nonlinear resonance regime and equal to $j = 1/3$ in the stochastic heating regime. The kinetic energy gain in the laboratory frame of reference $E_p$ can be found by performing the inverse Lorentz transformation from the boosted to the laboratory frame of reference,

$$E_p \sim \gamma M (E'_\text{total} - p_v V) - m_e c^2.$$ 

(29)

This is the equation for the proton energy $E_p$. Its analytical solution yields

$$E_p \sim \frac{m_e c^2}{\cos^2 \phi} \left[ \frac{2}{\pi} E \left( \frac{\phi}{1 + A} \right) \sqrt{1 + A} - 1 \right]$$

(30)

where

$$A = \frac{a_0^2 N}{\sin^2 \theta \sin^2 2\phi}$$

(31)

and

$$E(z) = \int_0^{\pi/2} \sqrt{1 - z^2 \sin^2 \Theta} d\Theta$$

(32)

is the complete elliptic integral of the second kind [31].

It should be emphasised that the present model determines the maximum energy of protons as a function of the

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Figure 3. Driven oscillations obtained for $a_0 = 2$ and $\varepsilon_p = 1.35$. Normalised momentum (a) and coordinate (b) of electrons versus time. (c) The phase plane $(x, p)$. (d) The Poincaré section showing the particle positions on the phase plane at the discrete time with the time step equal to the period of the driving force, $2\pi$. 

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(22)
laser parameters: dimensionless amplitude $a_0$, the number of wave period $N$, the angles of polarisation $\theta$ and incidence $\phi$. In figure 1, the analytically value of proton energy $E_p$ predicted with a nonlinear resonance mechanism (see equation (30) with $j = 1/2$) is shown as a function of $\theta$. Here, we are using the parameters employed in the experiment: $a_0 = 1.2$, $\phi = 45^\circ$, and $N = 15$, which corresponds to the number of wave periods included in the FWHM duration ($\tau = 40$ fs) of the incident laser. As a comparison, the proton energy analysed without considering the nonlinear effect in equation (30) ($j = 0$) is shown as a dashed line.

As is shown in equation (23), the present model is valid when the target thickness satisfies $l \simeq \Delta x$. The displacement of electrons $\Delta x$ is typically assumed to be equal to the amplitude of electron oscillation $r_E = eE_0/m_\omega^2$. Hence, the present model is applicable when the target thickness satisfies $l \simeq r_E$. To demonstrate applicability of the criterion above, we observed the target thickness dependency of proton energy, as shown in figure 4.

Figure 4 shows the dependence of the proton energy on the target thickness observed with $(a_0, \theta, \phi, N) = (1.2, 0, 45^\circ, 15)$. The maximum proton energy is close to the value of the analytical prediction (2.24 MeV) given by equation (30) when the target thickness ranges around $r_E = a_0\lambda 2\pi = 0.15 \mu m$. For thicker targets ($l \geq 5 \mu m$), the proton energy decreases with the increase in the target thickness $l$. This can be attributed to the neutralisation of the ion density $n_0$ by the returning current, whose effect is substantially stronger in the case of thicker targets.

4. Conclusion

In conclusion, We have studied the ion acceleration during the interaction of a short pulse laser obliquely incident on a thin foil target. The dependence of the ion energy on the laser polarisation has been studied experimentally and theoretically. We found that the ion energy being maximal for the $p$-polarisation gradually decreases when the pulse becomes $s$-polarised. The experimentally found dependencies of the ion energy are explained by invoking the anomalous electron heating, which results in efficient ion acceleration at the target surface. We showed that the proton energy can be substantially enhanced under the conditions of nonlinear resonance experienced by the electrons circulating around the thin foil irradiated by the electromagnetic wave. The nonlinear vacuum heating is efficiently realised when a high-contrast laser pulse is obliquely incident with the $p$-polarisation on a foil target, the thickness of which is appropriately chosen via a criterion of $l \approx r_E$. Under this condition, MeV-energy electrons and ions can be generated by a few terawatt femtosecond laser, leading to a compact laser-driven radiation source in the future.

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