High-energy monoenergetic proton beams from two stage acceleration with a slow laser pulse


Published in:
PHYSICAL REVIEW SPECIAL TOPICS-ACCELERATORS AND BEAMS

Document Version:
Publisher's PDF, also known as Version of record

Queen's University Belfast - Research Portal:
Link to publication record in Queen's University Belfast Research Portal

Publisher rights
© 2015 The Authors
Published by the American Physical Society under the terms of the Creative Commons Attribution 3.0 License (http://creativecommons.org/licenses/by/3.0/) Further distribution of this work must maintain attribution to the author(s) and the published article’s title, journal citation, and DOI.

General rights
Copyright for the publications made accessible via the Queen’s University Belfast Research Portal is retained by the author(s) and / or other copyright owners and it is a condition of accessing these publications that users recognise and abide by the legal requirements associated with these rights.

Take down policy
The Research Portal is Queen’s institutional repository that provides access to Queen's research output. Every effort has been made to ensure that content in the Research Portal does not infringe any person’s rights, or applicable UK laws. If you discover content in the Research Portal that you believe breaches copyright or violates any law, please contact openaccess@qub.ac.uk.
High-energy monoenergetic proton beams from two stage acceleration with a slow laser pulse

H. Y. Wang,1 X. Q. Yan,2,* and M. Zepf1,3,†

1Helmholtz Institute Jena, Fröbelstieg 3, 07743 Jena, Germany
2State Key Laboratory of Nuclear Physics and Technology, and Key Lab of High Energy Density Physics Simulation, CAPT, Peking University, Beijing 100871, China
3School of Mathematics and Physics, The Queen’s University of Belfast, BT7 1NN Belfast, United Kingdom

(Received 23 May 2014; published 3 February 2015)

We present a new regime to generate high-energy quasimonoenergetic proton beams in a “slow-pulse” regime, where the laser group velocity $v_g < c$ is reduced by an extended near-critical density plasma. In this regime, for properly matched laser intensity and group velocity, ions initially accelerated by the light sail (LS) mode can be further trapped and reflected by the snowplough potential generated by the laser in the near-critical density plasma. These two acceleration stages are connected by the onset of Rayleigh-Taylor-like (RT) instability. The usual ion energy spectrum broadening by RT instability is controlled and high quality proton beams can be generated. It is shown by multidimensional particle-in-cell simulation that quasimonoenergetic proton beams with energy up to hundreds of MeV can be generated at laser intensities of $10^{21}$ W/cm².

DOI: 10.1103/PhysRevSTAB.18.021302

PACS numbers: 52.38.Kd, 52.35.Mw, 52.59.-f

During the past decade, plasma-based ion accelerators have attracted a lot of attention due to their potential applications for particle acceleration, medical therapy [1], proton imaging, and inertial confinement fusion [2]. An important goal, developing sources of laser-driven protons for radiation therapy of deep-seated tumors [3], requires hundreds of MeV proton/ion beams with a small energy spread. Recent theoretical studies of the laser ion acceleration show that the radiation pressure acceleration (RPA) is a promising route to obtaining high-quality ion beams, including the hole boring (HB) regime for thick target [4–6] and light sail (LS) regime for thin target [7–16]. In particular, the LS scheme, where the ultra-thin target is accelerated by the radiation pressure as a whole, has drawn much interest due to the possibility to reach very high ion energies with foreseeable laser and target technology.

The simplest model of LS is analogous to the method proposed for the spacecraft propulsion [17], where the target plasma acts as an opaque plane mirror accelerated by the radiation pressure of the laser. The frequency of the reflected pulse due to the Doppler effect is $\omega_r = \omega(1-\beta)/(1+\beta)$, where $\omega$ is the incident pulse frequency and $\beta = V/c$ is the normalized velocity for the receding mirror. The conversion efficiency from the laser to the kinetic energy of the mirror is then given by: $\eta = 2\beta/(1+\beta)$. Hence, RPA becomes increasingly efficient ($\eta \to 1$) as $\beta \to 1$. Since the ions are accelerated to the same velocity as the electrons, most of the laser energy will be transferred to the ions due to their larger mass.

The predictions of the LS model are very appealing. However, most of the previous theoretical studies are based on one-dimensional simulations [9,14] or ultra-high intensity lasers $I > 10^{22} \text{ W/cm}^2$ with steep rising edge [7,12,13]. In more realistic multidimensional geometries and at lower intensities, the foil is very susceptible to the transverse Rayleigh-Taylor-like instability (RTI) [18,19]. It sets in at the very beginning of laser-foil interaction and causes rippling of the foil at the early linear stage. Gradually, the initial small perturbations grow exponentially with time and the foil quickly reaches the nonlinear stages of RTI, where the foil is torn into many high density spikes and bubbles [18]. During this process the ion energy spectrum is strongly broadened and finally often shows a quasiexponential decay with sharp cutoff energy, similar to the typical energy spectrum from target normal sheath acceleration [20].

In this article, we report on a new two stage acceleration regime of quasimonoenergetic proton beam generation, where a two layer target [solid density (SD) layer + near-critical density (NCD) layer] is illuminated by a circularly polarized laser pulse. In contrast to previously discussed schemes for LS, the laser pulse is greatly slowed...
down by the extended NCD plasma (laser group velocity \( v_g < c \)). Bulanov et al. has built the 1D ion acceleration model in such a slow wave RPA regime [21], which shows that the ion velocity cannot exceed the laser group velocity because the photons can no longer reach the receding mirror. We identify here that, however, this situation changes drastically in multidimensional case. In particular, we find that under proper coupling conditions between the laser intensity and laser group velocity, the development of RTI can act as a smooth connection for a second ion acceleration stage. The ions initially accelerated by LS can be further trapped and reflected by the slowly propagating laser snowplough potential in NCD plasma, and reach velocity higher than the laser group velocity. This provides a new route to controlling RTI which may otherwise prove detrimental to the LS scheme. 2D particle-in-cell (PIC) simulations show that high-density monoenergetic proton beams with a peak energy 300 MeV can be produced at moderate laser intensities of the order of \( 10^{21} \) W/cm².

For relativistic laser intensity, the NCD plasma becomes relativistically transparent if the plasma density is below a threshold density [22]. The group velocity, referred in the present paper as the on-axis laser front propagating velocity in NCD plasma, becomes slower than the light speed in vacuum. It is difficult to exactly estimate the group velocity of the laser pulse in NCD plasma because of both the nonlinearity in the wave equation and the self-consistent modification of the plasma density profile, and the situation is further complicated in multidimensional geometry due to transverse inhomogeneous laser radiation. By means of 2D simulations we find that the laser group velocity is almost unchanged over time in a uniform NCD plasma, and it decreases rapidly with the increase of plasma density for a fixed laser intensity. For such a slow laser pulse, the light pressure force can be expressed in terms of the flux of the EM wave momentum [7], which is proportional to the Poynting vector, \( \mathbf{S} = \mathbf{E} \times \mathbf{B}/4\pi \). Considering a circularly polarized electromagnetic wave propagating along the x axis with vector potential \( \mathbf{A} = A_0 [\cos(\omega t - kx) \mathbf{e}_y + \sin(\omega t - kx) \mathbf{e}_z] \), the Poynting vector is \( \mathbf{S} = \omega k c A_0^2 \mathbf{e}_z \). In the boosted frame of reference moving with the foil, we obtain [21],

\[
\mathbf{S}' = S_0 (\beta_g - \beta) (1 - \beta g \beta) / (1 - \beta^2).
\]

Here the primed values correspond to the moving reference frame, variables without a prime are taken in the laboratory frame. \( \beta \) and \( \beta_g \) are the foil velocity and laser group velocity in unit of \( c \), \( S_0 \) is the fixed Poynting vector for laser pulse in vacuum (\( \beta_g = 1 \)). Neglecting absorption for simplicity, the force acting on a thin foil is given by: \( \mathbf{F} = 2 R \mathbf{S}' \), where \( R \) is the reflectivity coefficient of the laser pulse. Then the equations of motion of the foil can be expressed as:}

\[
\frac{d}{dt} (\beta \gamma) = \frac{2I}{\rho' c^2} R (\beta_g - \beta) (1 - \beta g \beta) / (1 - \beta^2),
\]

where \( \gamma = 1 / \sqrt{1 - \beta^2} \), \( I \) is the laser intensity, \( \rho \) and \( \rho' \) are the mass density and thickness of the foil. Here we refer from Eq. (2) \( \beta_g = 1 \) to standard LS regime and \( \beta_g < 1 \) to the slow wave LS regime. Analytical solutions to Eq. (2) exist depending on suitable expressions for \( R \), while the simplest case being that of a constant \( R \). In the standard LS regime, the foil velocity can be obtained as a function of the pulse fluence \( \varepsilon \):

\[
\beta = \frac{(1 + \varepsilon)^2 - 1}{(1 + \varepsilon)^2 + 1},
\]

where \( \varepsilon = \frac{2R}{\rho' c^2} = \frac{2 R \pi^2 m e^2 z}{A_m c}, \quad a_0 = \frac{(0.85 / \sqrt{2}) \times (1 \lambda_0^2 / 10^{18} \text{ Wcm}^{-2} \text{ m}^2)^{1/2}}{\left( c / \lambda_0 \right)^{1/2}} \) is the dimensionless pulse amplitude for a circularly polarized laser pulse, \( \tau \) is the effective acceleration time in unit of laser period, and \( \zeta = \pi (n_0 / n_e) (\varepsilon / \lambda_0) \) which is the parameter governing the transparency of the thin solid density target [23]. In these equations, \( n_0 \) is the initial electron density, \( n_e = \pi m_e c^2 / e^2 \lambda_0^2 \) is the critical density, and \( \lambda_0 \) is the laser wavelength in vacuum. In the slow wave LS regime, Eq. (2) can only be solved in quadrature, it reads:

\[
\ln \frac{[1 - \beta \beta_g (1 - \beta)^{1/2} (1 - \beta^2)^{1/2} / \beta_g]}{(\beta_g - \beta) [1 + (1 - \beta_g^2)^{1/2}]} - \beta_g \left\{ \arctan \left( \frac{(1 - \beta_g^2)^{1/2} (1 - \beta^2)^{1/2}}{\beta_g - \beta} \right) - \arccos \beta_g \right\} = (1 - \beta_g^2)^{3/2} \varepsilon.
\]

In the limit of \( t \to \infty \) we have

\[
\beta = \beta_g - \exp\left\{ -(1 - \beta_g^2)^{3/2} \varepsilon \right\}.
\]

Thus the maximum ion velocity is limited by the laser group velocity. From Eq. (2) we can also find that at early time \( \beta \sim 0 \), the light pressure force \( F \propto \beta_g \), which means that the acceleration gradient is smaller in the slow wave LS regime. These arguments imply that ion acceleration in the slow wave LS regime is less efficient than the standard LS regime in 1D case.

However, as will be shown in 2D simulation results below, the acceleration dynamics become radically different in multidimensional geometry because several effects, including e.g., increased heating of electrons, transverse expansion of the foil, and onset of RTI, come into play. In the present paper we compare two cases: single solid density layer target and two layer target to explore the standard LS regime and slow wave LS regime, respectively. The 2D PIC simulations are run with the code KLA2D.
We choose a circularly polarized laser pulse with \( I = 2.47 \times 10^{21} \text{ Wcm}^{-2} \) \((\alpha_0 = 30, \lambda = 1.0 \mu \text{m})\) and super-Gaussian intensity distribution \( \exp[-(r/r_0)^4] \) propagating along the \( z \) axis, where \( r_0 = 6 \mu \text{m} \) is the spot radius. The laser pulse has a trapezoidal temporal profile, consisting of a plateau of \( 30T_0 \) and rising and falling times of \( 10T_0 \) each, where \( T_0 = 3.3 \) fs is the laser period. For the single layer target, a hydrogen layer with electron density \( n_0 = 100n_c \) and thickness 100 nm is located at \( z = 5 \mu \text{m} \), which satisfies the requirement of the optimal LS acceleration [9]. For the two layer target, an extended near-critical density layer, consisting of fully ionized carbons with thickness 30 \( \mu \text{m} \), is attached behind the hydrogen layer. The electron density in the carbon layer is set to \( 40n_c \) for critical density layer, consisting of fully ionized carbons with thickness 30 \( \mu \text{m} \), is attached behind the hydrogen layer. The whole acceleration process for the two cases is shown in Fig. 1. We find that at early stage for both cases, almost the whole laser-irradiated area is detached and pushed forward by the radiation pressure. Ions undergo almost the whole laser-irradiated area is detached and pushed forward by the radiation pressure. Ions undergo a typical LS stage, showing a clear peak in the energy spectrum. For the single SD foil, in contrast to 1D case a limited time due to the development of RTI, as seen in Fig. 1(b), (e), and (f). In our simulations, the characteristic time for RTI development \( \tau_{RT} \) is shorter than the pulse duration \( \tau_L \), and the effective acceleration time \( \tau = \tau_{RT} \). Using the linear stability theory, the RTI growth time follows \( \tau_{RT} \propto 1/\sqrt{k\rho} \) [25], where \( k = \frac{2\pi}{x_{RT}} \) is the instability
wave number and $g$ is the acceleration. On the one hand, the wavelength of the instability modulation is limited by diffractive effects [4] to

$$\lambda_{RT} \sim \lambda = \lambda_0 / \beta_g,$$  \hspace{1cm} (6)

which implies a correspondingly larger modulation wave length for a slower laser pulse. This can be clearly seen in Fig. 1(b), (c) as well, the modulation wavelength $\lambda_{RT}$ is $\sim \lambda_0$ for the single layer target and $\sim 1.8 \lambda_0$ for the two layer target, consistent with Eq. (6). On the other hand, in non-relativistic limit ($\beta \ll 1$), from Eq. (2) we find the acceleration $g \propto F \propto \beta_g$. Thus we can obtain the characteristic time for RTI

$$\tau_{RT} = \tau_0 / \beta_g,$$ \hspace{1cm} (7)

where $\tau_0$ is the characteristic time for RTI at $\beta_g = 1$. Following Eq. (7) one may expect a longer LS acceleration in the slow wave LS regime, which compensates for the decrease of acceleration gradient. This is also confirmed by simulation results seen in Fig. 2 (black and red stars), which show that the two layer target exhibits efficient LS for about twice the time.

Following the above arguments in 2D case one expects a longer and more stable ion LS acceleration from the two layer target. However, this stable LS will eventually be terminated when the foil reaches the nonlinear stage of RTI. As seen in Fig. 1(b), the electron layer from NCD plasma is greatly dispersed by RTI, showing several modulated density peaks. The proton distribution in space exhibits many clumps and bubbles, which is a very characteristic feature for the nonlinear RTI stage. The laser radiation selectively bores through the foil at the places where the foil becomes dilute, in analogy to fluid mixing in standard RTI. As shown in Fig. 1(d) at $t = 40T$, the proton energy spectrum is broadened and the energy peaks observed initially disappear. In fact, most previous studies in LS regime show a similar result [8,10,18,26]. The attractive features of the LS regime (i.e., high beam density and narrow energy spectrum) are destroyed at later time by the fast growth of RTI.

Surprisingly, as shown by the red dots in Fig. 2 and seen in Fig. 1(d), a new proton energy peak appears at later time (after $t = 40T$) for the two layer target. The ion peak velocity exceeds the laser group velocity, implying a new second ion acceleration stage after the initial LS stage. These two acceleration stages are connected by the development of the nonlinear RTI. As will be further described below, this new second acceleration stage is related to reflection of the LS ions by the slow propagating snowplough potential, similar to the hole boring and/or shock wave reflection process discussed in previous papers [4–6,27–29].

When the two layer target reaches the nonlinear RTI stage, the laser relativistically propagates in the NCD plasma with the group velocity $v_g$. One fundamental phenomenon of laser interaction with NCD plasma is the piling up of electrons at wave front by the laser ponderomotive force, referred to as a snowplough. The spatial charge separation between the carbon ions and electrons in the region depleted of electrons just before the snowplough sets up a propagating electrostatic potential $\phi_{sp}$ that follows the snowplough. This snowplough potential moves at the laser group velocity, and its magnitude can be estimated to equate the laser ponderomotive potential $e \phi_{sp} = (\sqrt{1 + g_0^2} - 1) m_e c^2$. From our simulation parameters, $e \phi_{sp} \approx 30 m_e c^2 \ll \frac{1}{2} m_p \frac{A}{Z} v_g^2 \sim 617 m_e c^2$, this means that the background carbon ions cannot be trapped by this snowplough potential. However, for the protons initially accelerated by the LS stage to velocity $v_1$, in non-relativistic limit they may satisfy

$$e \phi_{sp} > \frac{1}{2} m_p (v_g - v_1)^2.$$ \hspace{1cm} (8)

These protons may then be reflected by the snowplough potential to a velocity $v_2 = \frac{2v_g - v_1}{1 - \frac{2v_g - v_1}{v_g}} = v_g$, thus exceeding the laser group velocity $v_g$.

The two stage acceleration process can be clearly seen from the proton phase space evolution shown in Fig. 3. At early stages ($t = 24T$ and $t = 32T$), the protons show a clear “spiral structure” [8,14], which is a typical feature for the LS ion acceleration. This “spiral structure” disappears at $t = 40T$ and the proton beam debunches in the longitudinal direction, implying the termination of the LS stage by the development of RTI, as shown in Fig. 1. At later time ($t = 48T$ and $t = 56T$), on the one hand, the ions that
The color bar represents the proton numbers in arb.unit of log density distributions along z-axis (averaged over $|y| < 1 \mu m$). The color bar represents the proton numbers in arb.unit of log density distributions along z-axis (averaged over $|y| < 1 \mu m$).

satisfy Eq. (8) are trapped and reflected by the laser snowplough potential, forming a pronounced density peak (see $t = 56T$). Considering that the peak ion velocity at end of LS stage is $v_1 \sim 0.46c$ and the laser group velocity $v_g = 0.58c$, we can estimate a reflected ion velocity of $v_2 = 0.65c$, consistent with the simulation results in Fig. 3. On the other hand, some slow ions that do not satisfy Eq. (8) are dispersed in space, forming a long tail behind the reflected ions. Finally, a quasimotoenergetic proton beam with the peak energy $\sim 300$ MeV, density $5n_c$ and the divergence $< 10^\circ$ is obtained at $t = 56T$, as seen in Fig. 1 and Fig. 3. Note that in Fig. 1(d) that the proton energy spectra are almost the same at $t = 52T$ and $t = 60T$. This can be understood in terms of the increased velocity of the reflected ions separating them from the snowplough potential, and thus preventing further interaction with the laser pulse. Furthermore, the narrow spectral feature is stable as long as Coulomb self forces are negligible. Here the background electrons provide Debye shielding and the spectral shape remains stable while the bunch propagates in the NCD plasma. The natural expansion of the bunch due to the intrinsic beam divergence will ultimately reduce the self-fields to a level where the spectral modification due to self-fields becomes negligible. Thus one may expect that the high quality proton beam will persist in time to the end with a sufficiently thick NCD layer.

As shown by Eq. (8), the second ion acceleration stage requires that ions reach enough high velocity in the first LS stage $v_{th} \sim v_g - \frac{2(1+\alpha_0^2-1)}{4\alpha_0^2}c$. Put this $v_{th}$ into Eq. (2) and we can obtain the required time $\tau_1$ for LS stage. As discussed above, the LS stage is terminated by RTI and the effective acceleration time is limited by $\tau_{RT}$. Therefore the condition for the second ion acceleration stage is $\tau_1 < \tau_{RT}$.

Figure 4(a) shows the scaling of $\tau_1$ and $\tau_{RT}$ for different laser group velocities. Note that the condition $\tau_1 < \tau_{RT}$ is satisfied when the laser group velocity is slower than a threshold value $\beta_{th} = 0.62$. To verify this, simulations for varied NCD plasma density of $2n_c$, $4n_c$, $6n_c$, $8n_c$, and $10n_c$ were performed with corresponding laser group velocities of $0.7c$, $0.65c$, $0.58c$, $0.45c$, and $0.35c$, respectively. We can see in Fig. 4(a) that the characteristic times for RTI observed in simulations roughly follows the theory prediction by Eq. (7), and $\tau_1 < \tau_{RT}$ is satisfied for plasma density higher than $4n_c$ (the corresponding laser group velocity is a little higher than theory predicted $\beta_{th}$).

Figure 4(c) shows the proton energy spectrum at $t = 60T$ for varied laser group velocity. A clear energy peak is observed for $\beta_g < 0.65$ (plasma density $> 4n_c$), in good agreement with the theory prediction. In contrast for
the NCD plasma [24], which increases the laser intensity in the 2D case. We note that in 3D simulations the Gaussian pulse (which decreases the acceleration gradient) and development of RTI (which terminates the acceleration regime) act together in the present scheme and turn into a positive effect for generating the high-energy quasimonenergetic proton beams. The two layer target may be fabricated by depositing carbon nanotube foams (NCD layer) on thin diamond-like-carbon foils (SD layer) [30]. This regime might open a new way to high quality proton beam generation at moderate laser intensity in the near future.

X. Q. Yan is supported by National Basic Research Program of China (2013CBA01502), NSFC (11025523) and National Grand Instrument Project (2012YQ030142).


