Ion acceleration by radiation pressure in thin and thick targets

Andrea Macchi[a,b], Carlo Benedetti[c]

[a]CNR/INFM/polyLAB, Pisa, Italy
[b]Dipartimento di Fisica "Enrico Fermi", Largo Bruno Pontecorvo 3, I-56127 Pisa, Italy
[c]Dipartimento di Fisica, Università di Bologna and INFN, Via Irnerio 46, I-40126, Bologna, Italy

Abstract

Radiation Pressure Acceleration (RPA) by circularly polarized laser pulses is emerging as a promising way to obtain efficient acceleration of ions. We briefly review theoretical work on the topic, aiming at characterizing suitable experimental scenarios. We discuss the two reference cases of RPA, namely the thick target ("Hole Boring") and the (ultra)thin target ("Light Sail") regimes. The different scaling laws of the two regimes, the related experimental challenges and their suitability for foreseen applications are discussed.

Key words: Laser-plasma interaction, ion acceleration, radiation pressure, circular polarization
PACS: 52.38.Kd, 41.75.Jv, 52.50.Jm

1. Introduction

In the interaction of superintense ($I = 10^{17} - 10^{22}$ W cm^{-2}) laser pulses with solid targets, two general mechanisms can lead to the acceleration of large numbers of ions to high energy (> MeV). The first mechanism is related to the generation of multi–MeV electrons by nonlinear laser-plasma interaction processes. Such electrons, by attempting to leave the target and escape in vacuum, lead to the generation of very intense space-charge fields which drive the acceleration of ions. This process is the basis of the Target-Normal Sheath Acceleration mechanism which accounts for most of the experimental observations of proton and heavier ion acceleration since the year 2000 (see [1] for a review).

The second mechanism is related to the radiation pressure of the laser pulse on the target. If the laser pulse is modeled as a plane wave and the target as a semi-infinite, opaque object of reflectivity R, the radiation pressure $P_{\rm rad}$, i.e. the momentum transferred by the wave to the target per unit time and surface, at normal incidence and in the rest frame of the target is given by the well-known expression due to J. C. Maxwell

$$P_{\rm rad} = (1+R)\frac{I}{c}. \qquad (1)$$

From a local and "microscopic" point of view, $P_{\rm rad}$ is given by the integral over the target volume of the temporal average of the $\mathbf{J} \times \mathbf{B}$ force over a period of the electromagnetic (EM) field, i.e. of the ponderomotive force per unit volume $\mathbf{f}_p$. Since $\mathbf{f}_p$ scales with the inverse of the particle mass, its "direct" effect on ions is negligible; the transfer of momentum to ions is actually mediated by the electrostatic field generated by the displacement of the electrons under the action of $\mathbf{f}_p$. If $\mathbf{f}_p$ and the electrostatic force on electrons balance each other ($\mathbf{f}_p - en_e\mathbf{E}_{\rm es} \simeq 0$), so that electrons are in mechanical equilibrium, the total electrostatic pressure $P_{\rm es}$ on ions is given by the integral of $Zen_i\mathbf{E}_{\rm es}$. If the conditions of quasi-neutrality $Zen_i = en_e$ holds, then $Zen_i\mathbf{E}_{\rm es} = en_e\mathbf{E}_{\rm es} = \mathbf{f}_p$, yielding $P_{\rm es} = P_{\rm rad}$. In these conditions one can effectively assume that the target ions are pushed by $P_{\rm rad}$.

In general, both TNSA and RPA are at play in laser-solid target interaction, but there is plenty of evidence that TNSA dominates in most of experiments reported so far on ion acceleration from solid targets. Only in a few recent publications [2, 3] significant radiation pressure effects were observed at very high intensities, in qualitative agreement with the trend observed in simulations which suggest RPA to become progressively more important at higher intensities and finally to dominate over TNSA at intensities $I > 10^{23}$ W cm^{-2} [4, 5], still beyond present-day laser systems capabilities. However, these considerations apply to linearly

Email addresses: macchi@df.unipi.it (Andrea Macchi), Carlo.Benedetti@bo.infn.it (Carlo Benedetti)
URL: http://www.df.unipi.it/~macchi (Andrea Macchi)

polarized pulses only. In fact, for circularly polarized pulses and normal incidence, the generation of energetic electrons is strongly suppressed ruling TNSA out, so that RPA dominates at any intensity (strictly speaking, at any intensity for which $P_{\rm rad}$ is larger than the thermal pressure related to collisional heating of the target). A large number of recent theoretical papers has investigated RPA driven by circularly polarized pulses. Based on these papers, two different regimes of RPA can be distinguished by the value of the target thickness. The first regime is that of "thick" targets, also known as the "hole boring" (HB) regime, where only ions in a surface layer of the target are accelerated by RPA [6, 7, 8, 9, 10, 11, 12]. The second regime, also named "Light Sail" (LS), is that of targets thin enough to be accelerated as a whole [13, 14, 15, 16, 17, 18, 19, 20, 21, 22]. In the following we characterize the HB and LS regimes summarizing for both the most relevant results and the relevant scaling laws, and commenting on the related experimental challenges and expected suitability for foreseen applications. While in the present paper we consider only basic features for briefity, we note that other recent papers addressed other related issues such as RPA in structured or composite targets [23, 24, 25, 26] or the effect of elliptical (rather than perfecly circular) polarization [27, 28]. Finally, we notice that first preliminary experimental results on RPA with circular polarization were communicated very recently, and in particular at the workshop to which the present journal issue is dedicated [29, 30, 31].

2. Thick targets

A model of ion acceleration in semi-infinite targets with homogeneous (step-like) density profile was proposed in Ref.[6]. The model is best described with the help of the cartoon in Fig.1. The first frame, Fig.1 a) corresponds to the initial stage in which the electrons have piled up under the action of the radiation pressure, creating the space-charge field E_x which balances the ponderomotive force. Assuming that the electrons are in equilibrium and ions have not moved significantly yet at this stage, the simple profiles of E_x and of the electron density n_e can be considered as simple approximations to the exact profiles that can be calculated in steady conditions [32]. The model parameters x_d, E_0 and n_{p0} (shown in Fig.1) are related to each other by the Poisson equation, the constraint of charge conservation and the equilibrium condition

$$P_{\rm es} = \int e n_e E_x dx \doteq P_{\rm rad} = 2\frac{I}{c}, \tag{2}$$

assuming total reflection ($R = 1$). The distance $\ell_s = x_s - x_d$ is the penetration distance of the ponderomotive force into the target, and is thus of the order of the collisionless skin depth $d_p = c/\omega_p$, where ω_p is the plasma frequency. The simplified profiles allow the ion equation of motion to be solved analytically, and it is easily found that all ions initially in the $x_d < x < x_s$ region get to the $x = x_s$ point at the same time τ_b, so that a singularity appears in the ion density n_i. The corresponding energy spectrum is a flat-top distribution extending from zero to the cut–off value (per nucleon)

$$E_{\rm max} = \frac{m_p}{2}\upsilon_m^2 = 2m_p c^2\Pi, \tag{3}$$

where

$$\frac{\upsilon_m}{c} = 2\Pi^{1/2}, \qquad \Pi = \frac{I}{m_i n_i c^3} = \frac{Z}{A}\frac{n_c}{n_e}\frac{m_e}{m_p}a_0^2. \tag{4}$$

In these equations, $n_c = 1.1\times10^{21}\ \mathrm{cm}^{-3}(\lambda/1\ \mu\mathrm{m})^{-2}$ is the cut–off density corresponding to the laser wavelength λ and $a_0 = (0.85/\sqrt{2})(I\lambda^2/10^{18}\ \mathrm{W\ cm}^{-2}\mu\mathrm{m}^2)^{1/2}$ is the dimensionless pulse amplitude corresponding to the intensity I (for circular polarization). Actually, the "final" energy spectrum observed in the simulations is determined by the highly transient "wave–breaking" stage following the formation of the singularity (corresponding to a sharp density spike). The fastest ions form a narrow bunch of velocity υ_m that penetrates into the overdense plasma. The other ions form a second peak moving at approximately the average velocity $\upsilon_b = \upsilon_m/2$. Thus, υ_b is equal to the recession velocity of the radiation pressure–driven surface of the plasma whose expression can also be obtained using an argument of momentum flux balance between the laser pulse and the ion current [33, 10]. In other terms, υ_b is the speed at which a hole is bored in the plasma by the laser pulse, so that it is clear why this regime was also named as "hole boring".

An extension of Eq.(3) for the case of relativistic ion velocities was proposed in Refs.[10, 12]. The central point is that the momentum balance equation must be used in the reference frame comoving with the surface, where the intensity is $I' = I(1-\upsilon_b/c)/(1+\upsilon_b/c)$. One thus obtains

$$E_{\rm max} = 2m_p c^2\frac{\Pi}{1+2\Pi^{1/2}}. \tag{5}$$

Note that this approach accounts for pump depletion effects during the acceleration. In fact, assuming that the acceleration process is adiabatic, the "number of photons" is conserved in the reflection of the EM wave from the plasma surface and is a relativistic invariant. Thus,

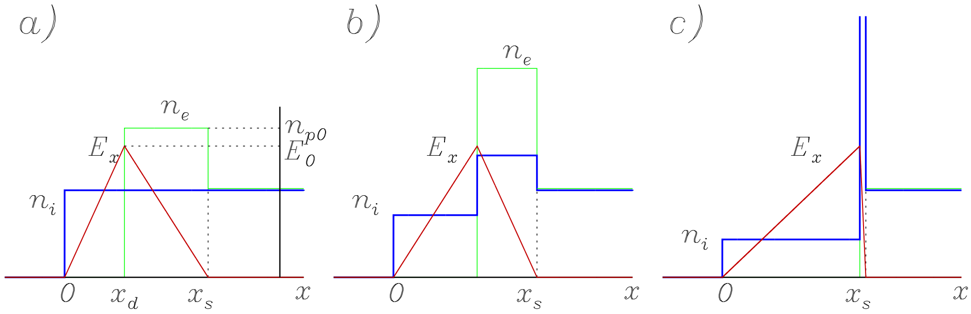

Figure 1: (Color online) Model profiles [6] of the ion density n_i (blue), electron density n_e (green) and electrostatic field E_x (red) at three different stages of ion acceleration.

the transfer of energy from the EM wave to the plasma in the laboratory frame occurs because of the frequency downshift of each photon that is reflected from the moving surface of the plasma. The fraction of the energy of the laser pulse which is transformed into kinetic energy is proportional to the relative frequency shift of the reflected wave, $\omega'/\omega = (1 - \upsilon_b/c)/(1 + \upsilon_b/c)$ which is the factor multiplying I in the "relativistic" case. By using the same argument one also obtains a simple and general formula for the instantaneous efficiency η, i.e. for the ratio between the mechanical energy delivered to the target and the incident pulse energy:

$$\eta = \frac{2\beta}{1+\beta}, \tag{6}$$

where $\beta = \upsilon_b/c$ for the HB regime.

The scaling of the ion energy with the laser intensity and the target density predicted by Eqs.(3-5) is unfavorable if one aims at relativistic ion energies. For targets at solid density ($n_e > 100n_c$), values of $\Pi \sim 1$ cannot be approached even at the highest intensities available today. Already to obtain $\Pi \sim 10^{-1}$, that would correspond to ions of energy $\sim 10^2$ MeV per nucleon which could be appropriate for applications such as hadrontherapy, target with density of only a few times n_c would be needed. A simple strategy to obtain such densities might be to preform a short-scalelength plasma on the target surface, e.g. by a short prepulse or pedestal preceding the main ultrashort pulse. Some preliminary simulation results showed that in inhomogeneous plasmas with a scalelength of a few wavelengths at the $n_e = n_c$ surface ion acceleration occurs very similarly to the case of a step-like profile, and that the observed ion energies imply that the relevant value of the electron density, as given by Eq.(3), is just a few times n_c [18].

A regime of HB-RPA producing ions with energy exceeding 10^2 MeV might be obtained with the foreseen development of few-cycle pulses at ultrahigh intensity. The extremely short duration, possibly shorter than the bunch formation time τ_b, would allow to concentrate all the pulse energy in the acceleration of a single ion bunch. Moreover, if efficient conversion to circular polarization can be achieved for such "extreme" laser pulses, it might allow to make ion acceleration stable with respect to the variation of the "absolute" phase, because the slowly varying ponderomotive force depends only on the pulse envelope.

Fig.2 shows spectra from two-dimensional (2D) particle-in-cell (PIC) simulations of the interaction of a two-cycle duration, circularly polarized pulse with a "liquid hydrogen" target having $n_e = 50n_c$. The peak amplitude of the pulse is $a_0 = 108$ (corresponding to $I\lambda^2 = 3.2 \times 10^{22}$ W/cm$^2\mu$m^2), and the pulse intensity has a longitudinal profile given by a $\sin^4$ function with a FWHM of 1.45λ and a Gaussian transverse profile at

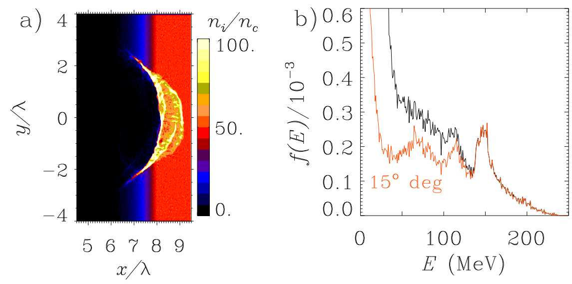

Figure 2: (Color online) 2D simulation of the interaction of a two-cycle pulse with a "liquid hydrogen" slab. a): density contours after the pulse peak. b): the corresponding energy spectrum. The lower (red) line is the spectrum restricted to ions having velocity into a cone of 15^o aperture around the axis. See text for parameters.

the waist with a FWHM of 2.2λ. The plasma density increases as a $\sim x^4$ function over a length of 4λ and then remains constant. As shown in Fig.2 a), a bunch of ions is generated at the front surface. The spectrum in Fig.2 b) shows a peak at an energy of about 150 MeV, corresponding to ions travelling near the axis; ions of lower energy are accelerated at the edges of the laser spot where the intensity is lower, as can be inferred by the spectrum restricted to ions directed within a 15 degrees cone centered on axis. The 2D simulations were performed with the `ALaDyn` code [34].

A liquid hydrogen jet might be an interesting option for HB-RPA with ultraintense few-cycle pulses because, besides allowing the acceleration of protons, being a "continously flowing" target it would allow for high repetition rate operation, which is of paramount importance for applications. Using gas jets as flowing targets for HB-RPA might work with long wavelength lasers, e.g. CO_2, for which the electron density can be varied in the range up to a few times n_c .

3. Thin targets

For long enough values of the pulse duration τ_L such that $\upsilon_b\tau_L > \ell$, being ℓ the thickness of the target, the hole boring through the target is complete and all the ions are displaced by the radiation pressure action. For a very thin target having $\ell \lesssim \ell_s$ all the ions in the foil are accelerated as a single bunch, and the acceleration cycle can be repeated for the whole duration of the laser pulse [13, 15, 16] (we remind that $\ell_s = x_s - x_d$ in Fig.1 is the effective evanescence length of the ponderomotive force). The average motion of the target can be described as that of a reflecting object (i.e. a mirror) boosted by the radiation pressure, hence the definition of "Light Sail" regime. The appealing features of LS are the monoenergetic ion spectrum expected for a "rigid" displacement of the whole foil target and the possibility to achieve high energies due to repeated acceleration and the low surface mass. However, for such thin targets the transmission of the laser pulse through the target must be taken into account. The radiation pressure on a thin target in its rest frame, at normal incidence and neglecting absorption, is given by

$$P_{\rm thin} = 2R_{\rm thin}\frac{I}{c} \tag{7}$$

where $R_{\rm thin}$ is the reflectivity of the thin target. Eq.(7) is different from the expression for a thick target, Eq.(1), because the trasmitted wave must be taken into account. Under the action of $P_{\rm thin}$, the foil velocity $V = \beta c$ satisfies the equation of motion

$$\frac{d}{dt}(\beta\gamma) = \frac{2I(t_{\rm rit})}{m_i n_i \ell c}\frac{1-\beta}{1+\beta}R_{\rm thin}(\omega'), \tag{8}$$

where $t_{\rm rit}$ is the retarded time at the foil position and $\omega' = \omega\sqrt{(1-\beta)/(1+\beta)}$ is the frequency in the rest frame of the foil.

For "relativistic" laser intensities, the reflectivity $R_{\rm thin}$ depends nonlinearly on the laser intensity because of self-induced transparency effects. The relevant parameter is $\zeta = \pi(n_e/n_c)(\ell/\lambda)$ [35, 22]. For $\zeta \gg 1$, $R_{\rm thin} \simeq 1$ for $a_0 > \zeta$ and drops abruptly as a_0^{-2} for $a_0 < \zeta$. Thus, the condition $a_0 = \zeta$ is the best compromise between increasing the boost on the foil and decreasing its mass. The existence of an "optimal" thickness for the foil acceleration is further discussed in various references, e.g. [17, 20, 19, 36].

For $R_{\rm thin} = 1$, which is a good approximation for a thin solid foil unless the thickness is just a few nanometers, Eq.(8) gives a remarkably simple result for the final velocity of the foil as a function of the pulse fluence, i.e. of the total energy incident for unit surface:

$$\beta_{\rm f} = \frac{(1+\mathcal{E})^2-1}{(1+\mathcal{E})^2+1}, \qquad \mathcal{E} = \frac{2}{\rho\ell c^2}\int I(t')dt' \tag{9}$$

where in terms of the pulse and foil parameters

$$\mathcal{E} \simeq 2\pi\frac{Z}{A}\frac{m_e}{m_p}\frac{a_0^2\tau_L}{\zeta} \tag{10}$$

where τ_L is now in units of the laser period. A similar formula for a foil with $R_{\rm thin} < 1$, in which $\beta_{\rm f} = \beta_{\rm f}(\mathcal{E},\zeta)$ has also been obtained recently [36].

The efficiency is again the same function of the foil velocity as given in Eq.(6), because the same argument applies. Thus, the higher efficiency predicted for LS with respect to HB is merely a consequence of the possibility to reach higher values of β with feasible laser and target technology.

The scaling of the energy per nucleon $E/A = m_p c^2(\gamma_{\rm f} - 1)$ (with $\gamma_{\rm f} = (1-\beta_{\rm f}^2)^{-1/2}$) with the dimensionless fluence $a_0^2\tau_L$ is shown in Fig.3 a). Present–day laser systems may approach fluence values close to $a_0^2\tau_L = 10^4$ (corresponding, e.g., to a 30 fs laser pulse with an intensity of a few times 10^{21} W cm^{-2}) while foils as thin as a few nanometers, having $\zeta \sim 1\div10^1$ may be produced. In such conditions, according to Fig.3 a) the GeV barrier may be approached in the LS regime.

PIC simulations of thin foil acceleration in one dimension (1D) show that the ion spectrum is dominated by a single spectral peak for intensities such that $a_0 < \zeta$,

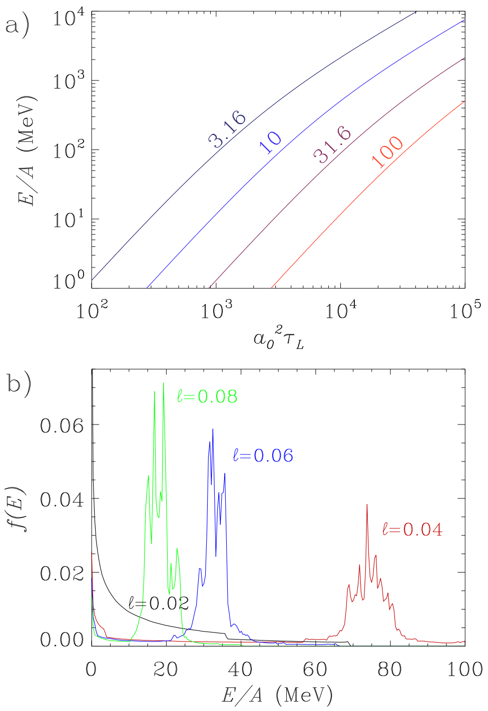

Figure 3: (Color online) a): Energy scaling given by the LS model for $R = 1$. The energy per nucleon is shown as a function of the dimensionless fluence of the laser pulse $a_0^2\tau_L$ (with τ_L the duration in units of the pulse period) and of the foil parameter ζ. b): Energy spectra from 1D PIC simulations of the interaction of a laser pulse with plasma foils of different thickness ℓ (in units of λ). Parameters common to all simulations are $a_0 = 30$, $n_e/n_c = 250$, $\tau_L = 8$, $Z/A = 1/2$.

as shown in Fig.3 b). For $a_0 > \zeta$, most of the electrons are pushed out of the foil which explodes due to Coulomb repulsion, and an exponential–like spectrum is observed as for the $\ell = 0.02\lambda$ case in Fig.3 b). The peak energy is in very good agreement with the predictions of LS model. At the same time, however, the spectral peak contains only a fraction of the target ions, corresponding approximately to the ions starting from the $x_d < x < x_s$ region in Fig.1, which may be much thinner than the foil thickness. This sounds as a paradox, because if just a fraction of the foil is accelerated, its inertia should be lower than that of the original foil and lead to higher velocity. An explanation has been given in Ref.[22]: due to the charge unbalance in the $x_d < x < x_s$ region, so that the accelerated layer is negatively charged, the electrostatic pressure $P_{\rm es}$ is lower than the radiation pressure $P_{\rm thin}$. By using the profiles of Fig.1 it can be shown than the ratio $P_{\rm es}/P_{\rm thin} = \ell_{\rm acc}/\ell$, where $\ell_{\rm acc}$ is the thickness of the accelerated layer; thus, in the equation of motion of the layer the two reduction factors in the foil mass and in the boosting pressure cancel each other, so that the equation of motion effectively preserves the form of Eq.(8) and has the same solution.

Numerical simulations shows that while the simple LS model of the “accelerating mirror” is successful in predicting the ion energy, the underlying dynamics is much more complex (already in 1D) and shows a high degree of self-organization [22, 20, 21, 36]. Such dynamics, of which the above described effect is a prominent example, originates in the different action of the ponderomotive force of electrons and ions, driving strong charge separation. Another important effect is the heating of electrons near the end of the laser pulse, when the radiation pressure drops down and cannot maintain the force equilibrium. Heating of electrons may cause a significant post-acceleration expansion of the accelerated layer leading to broadening of the ion spectra, as in Fig.3 b). Thus, already in 1D geometry, the LS regime may not lead to a monoenergetic spectrum as desired.

Additional effects and issues are found in multi-dimensional simulations. The intensity distribution of the laser pulse in the focal plane is in general not uniform, causing the ion energy spectrum to broaden and becoming dependent on the angle with respect to the axis. Both the ion spectrum and collimation are improved by using “flat-top” intensity distributions (e.g. supergaussian profiles) [15, 16], which also prevent early target disruption and pulse transmission due to the expansion of the target in lateral direction, as found in 3D simulations [18]. However, even for flat-top intensity profiles effects such as target bending (leading to local oblique incidence) or surface instabilities [16, 8], causing electron heating and broadening of the ion spectrum. To reduce such detrimental effects, complex intensity distributions [8] or target with shaped density profiles [23] have been recently investigated.

4. Conclusions and outlook

There has been much theoretical and simulation work devoted to Radiation Pressure Acceleration both in the Hole Boring (thick target) and in the Light Sail (thin target) regimes. Circularly Polarized pulses have been considered in almost all works to maximize the efficiency of RPA. The concept is appealing for ion acceleration, but a number of open and challenging issues, such as obtaining a monoenergetic spectrum, is apparent already at the level of low-dimensional simulations. As first experimental results on RPA are now appearing, an even more intense activity oriented to the interpretation and modeling of experimental results is expected.

The scaling of ion energy with the laser intensity in makes difficult to obtain high ion energies in the HB regime, which may be appropriate indeed for applications requiring large number of ions with energies up to a few MeV, such as heating of high-density matter. Recently the HB regime was proposed as an option for ion-driven Fast Ignition in Inertial Confinement Fusion [11]. In addition, the HB regime seems to be less prone to prepulse effects and thus be easier to observe in experiments. It also might be more suitable for high-repetition rate operation, if flowing targets of appropriate density can be developed. As we have shown with a simulation example, next-generation few-cycle pulses with intensities approaching 10^{23} W cm^{-2} may allow to reach energies exceeding 100 MeV also in the HB regime.

For the LS regime, the scaling with the pulse energy may allow to reach ion energies up to the GeV range, of interest for high-energy physics experiments, using near-future laser technology. The requirement of ultrathin (few nanometers) target makes the use of pulse cleaning techniques, such as plasma mirror, essential. For obvious reasons of target manufacturing it is more appropriate for the acceleration of ions with $Z > 1$, most likely Carbon ions from Diamond-Like Carbon foils.

For both the HB and the LS regimes, most of the theoretical and simulation work have been and will be devoted to identify schemes and strategies, such as laser pulse shaping or target engineering, to improve the quality of RPA-accelerated ion beams. At the same time, the high degree of self-organization evidence by simulations makes RPA a topic of great interest for basic laser-plasma physics, showing several open theoretical issues. These latter may include, in our opinion, the detailed nature the transition to the radiation pressure-dominated regime for linearly polarized pulses [4, 5], the modeling of ultrahigh-intensity regimes where ions become relativistic within a single laser cycle [4, 19], the origin of surface instabilities [37, 16, 8], and the absorption of the angular momentum of circularly polarized pulses [18].

Acknowledgments

Support from CNR via a RSTL project and use of supercomputing facilities at CINECA (Bologna, Italy) sponsored by CNR/INFM are acknowledged.

References

[1] M. Borghesi, J. Fuchs, S. V. Bulanov, A. J. MacKinnon, P. K. Patel, M. Roth, Fast ion generation by high-intensity laser irradiation of solid targets and applications, Fus. Sci. Techn. 49 (2006) 412.

[2] S. Kar, et al., Plasma jets driven by ultraintense-laser interaction with thin foils, Phys. Rev. Lett. 100 (2008) 225004.

[3] K. U. Akli, et al., Laser heating of solid matter by light-pressure-driven shocks at ultrarelativistic intensities, Phys. Rev. Lett. 100 (2008) 165002.

[4] T. Esirkepov, M. Borghesi, S. V. Bulanov, G. Mourou, T. Tajima, Highly efficient relativistic-ion generation in the laser-piston regime, Phys. Rev. Lett. 92 (2004) 175003.

[5] T. Esirkepov, M. Yamagiwa, T. Tajima, Laser ion-acceleration scaling laws seen in multiparametric particle-in-cell simulations, Phys. Rev. Lett. 96 (2006) 105001.

[6] A. Macchi, F. Cattani, T. V. Liseykina, F. Cornolti, Laser acceleration of ion bunches at the front surface of overdense plasmas, Phys. Rev. Lett. 94 (2005) 165003.

[7] T. V. Liseikina, A. Macchi, Features of ion acceleration by circularly polarized laser pulses, Appl. Phys. Lett. 91 (2007) 171502.

[8] M. Chen, A. Pukhov, Z. M. Sheng, X. Q. Yan, Laser mode effects on the ion acceleration during circularly polarized laser pulse interaction with foil targets, Phys. Plasmas 15 (2008) 113103.

[9] Y. Yin, W. Yu, M. Y. Yu, A. Lei, X. Yang, H. Xu, V. K. Senecha, Influence of target thickness on the generation of high-density ion bunches by ultrashort circularly polarized laser pulses, Phys. Plasmas 15 (2008) 093106.

[10] A. P. L. Robinson, P. Gibbon, M. Zepf, S. Kar, R. G. Evans, C. Bellei, Relativistically correct hole-boring and ion acceleration by circularly polarized laser pulses, Plasma Phys. Contr. Fusion 51 (2009) 024004.

[11] N. Naumova, T. Schlegel, V. T. Tikhonchuk, C. Labaune, I. V. Sokolov, G. Mourou, Hole boring in a DT pellet and fast-ion ignition with ultraintense laser pulses, Phys. Rev. Lett. 102 (2009) 025002.

[12] T. Schlegel, N. Naumova, V. T. Tikhonchuk, C. Labaune, I. V. Sokolov, G. Mourou, Relativistic laser piston model: Ponderomotive ion acceleration in dense plasmas using ultraintense laser pulses, Phys. Plasmas 16 (2009) 083103.

[13] X. Zhang, B. Shen, X. Li, Z. Jin, F. Wang, Multistaged acceleration of ions by circularly polarized laser pulse: Monoenergetic ion beam generation, Phys. Plasmas 14 (2007) 073101.

[14] X. Zhang, B. Shen, X. Li, Z. Jin, F. Wang, M. Wen, Efficient GeV ion generation by ultraintense circularly polarized laser pulse, Phys. Plasmas 14 (2007) 123108.

[15] A. P. L. Robinson, M. Zepf, S. Kar, R. G. Evans, C. Bellei, Radiation pressure acceleration of thin foils with circularly polarized laser pulses, New J. Phys. 10 (2008) 013021.

[16] O. Klimo, J. Psikal, J. Limpouch, V. T. Tikhonchuk, Monoenergetic ion beams from ultrathin foils irradiated by ultrahigh-contrast circularly polarized laser pulses, Phys. Rev. ST Accel. Beams 11 (2008) 031301.

[17] X. Q. Yan, C. Lin, Z. M. Sheng, Z. Y. Guo, B. C. Liu, Y. R. Lu, J. X. Fang, J. E. Chen, Generating high-current monoenergetic proton beams by a circularly polarized laser pulse in the phase-stable acceleration regime, Phys. Rev. Lett. 100 (2008) 135003.

[18] T. V. Liseykina, M. Borghesi, A. Macchi, S. Tuveri, Radiation pressure acceleration by ultraintense laser pulses, Plasma Phys. Contr. Fusion 50 (2008) 124033.

[19] B. Qiao, M. Zepf, M. Borghesi, M. Geissler, Stable GeV ion-beam acceleration from thin foils by circularly polarized laser pulses, Phys. Rev. Lett. 102 (2009) 145002.

[20] V. K. Tripathi, C. S. Liu, X. Shao, B. Eliasson, R. Z. Sagdeev, Laser acceleration of monoenergetic protons in a self-organized double layer from thin foil, Plasma Phys. Contr. Fusion 51 (2009) 024014.

[21] B. Eliasson, C. S. Liu, X. Shao, R. Z. Sagdeev, P. K. Shukla, Laser acceleration of monoenergetic protons via a double layer

emerging from an ultra-thin foil, New J. Phys. 11 (2009) 073006.
[22] A. Macchi, S. Veghini, F. Pegoraro, “Light Sail” acceleration reexamined, Phys. Rev. Lett. 103 (2009) 085003.
[23] M. Chen, A. Pukhov, T. P. Yu, Z. M. Sheng, Enhanced collimated GeV monoenergetic ion acceleration from a shaped foil target irradiated by a circularly polarized laser pulse, Phys. Rev. Lett. 103 (2009) 024801.
[24] X. Zhang, B. Shen, L. Ji, F. Wang, Z. Jin, X. Li, M. Wen, J. R. Cary, Ion acceleration with mixed solid targets interacting with circularly polarized lasers, Phys. Rev. ST Accel. Beams 12 (2009) 021301.
[25] A. P. L. Robinson, D.-H. Kwon, K. Lancaster, Hole-boring radiation pressure acceleration with two ion species, Plasma Phys. Contr. Fusion 51 (2009) 095006.
[26] M. Grech, S. Skupin, R. Nuter, L. Gremillet, E. Lefebvre, High-quality ion beams by irradiating a nano-structured target with a petawatt laser pulse, New J. Phys. 11 (2009) at press, e–print arXiv:0906.3972.
[27] S. G. Rykovanov, J. Schreiber, J. Meyer–ter–Vehn, C. Bellei, A. Henig, H. C. Wu, M. Geissler, Ion acceleration with ultra-thin foils using elliptically polarized laser pulses, New J. Phys. 10 (2008) 113005.
[28] A. Macchi, T. V. Liseikina, S. Tuveri, S. Veghini, Theory and simulation of ion acceleration with circularly polarized laser pulses, C. R. Physique 10 (2009) 207 – 215.
[29] M. Borghesi, talk at the Coulomb09 Workshop – Ions Acceleration with high Power Lasers: Physics and Applications (Senigallia, Italy, June 15-18, 2009).
[30] X. Yan, talk at the Coulomb09 Workshop – Ions Acceleration with high Power Lasers: Physics and Applications (Senigallia, Italy, June 15-18, 2009).
[31] A. Henig, et al., Radiation pressure acceleration of ion beams driven by circularly polarized laser pulses, e-print arXiv:0908.4057 (2009).
[32] F. Cattani, A. Kim, D. Anderson, M. Lisak, Threshold of induced transparency in the relativistic interaction of an electromagnetic wave with overdense plasmas, Phys. Rev. E 62 (2000) 1234–1237.
[33] S. C. Wilks, W. L. Kruer, M. Tabak, A. B. Langdon, Absorption of ultra-intense laser pulses, Phys. Rev. Lett. 69 (1992) 1383–1386.
[34] C. Benedetti, A. Sgattoni, G. Turchetti, P. Londrillo, `ALaDyn` : A High-Accuracy PIC Code for the Maxwell-Vlasov Equations, IEEE Trans. Plasma Science 36 (2008) 1790–1798.
[35] V. A. Vshivkov, N. M. Naumova, F. Pegoraro, S. V. Bulanov, Nonlinear electrodynamics of the interaction of ultra-intense laser pulses with a thin foil, Phys. Plasmas 5 (1998) 2727–2741.
[36] A. Macchi, S. Veghini, T. V. Liseykina, F. Pegoraro, Radiation pressure acceleration of ultrathin foils, New J. Phys.Submitted.
[37] F. Pegoraro, S. V. Bulanov, Photon bubbles and ion acceleration in a plasma dominated by the radiation pressure of an electromagnetic pulse, Phys. Rev. Lett. 99 (2007) 065002.