Laser-driven generation of collimated quasi-monoenergetic proton beam using double-layer target with interface modulations

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Usage of double-layer targets consisting of heavy and light material with modulated interface between them provides a way for laser-driven generation of collimated ion beams. With extensive 2D3V PIC simulations we show that this configuration may result in a development of a relativistic instability with Rayleigh-Taylor and Richtmyer-Meshkov like features. Initially small perturbations are amplified during the laser-target interaction leading to the formation of low-density plasma regions and high-density bunches between them, which are accelerated by the laser radiation pressure as whole compact structures. That results in collimated quasi-monoenergetic proton beam with high average energy. The properties of this proton beam such as its low emittance (one order of magnitude lower compared to that of conventional accelerators) and divergence are discussed. Results are compared with similar acceleration schemes such as double-layer target without corrugation and single-layer target.

I. INTRODUCTION

Laser-driven ion accelerators has received a great deal of interest in last several decades as they are capable of sustaining relatively higher accelerating gradients than their conventional counterparts, and are currently able to accelerate protons to energies of 100 MeV [1]. With the advent of multi-petawatt laser systems like ELI-Beamlines (Czech Republic), APOLLON (France) or SEL (China) the laser pulses will soon reach intensities over $10^{23}$ W/cm$^2$ and will enter the acceleration regimes dominated by radiation pressure [2, 3] promising proton/ion acceleration above several GeV.

Studies of these high-intensity laser interaction with single-layer planar targets [4] shows the development of relativistic Rayleigh-Taylor [5, 6] like instability (RTI) leading to the formation of low density regions and high density ion bunches between them. The bunches exhibit quasi-monoenergetic behavior [4, 7]. The instability can develop in a controlled way, when a corrugation is imprinted on the front surface of the target [8]. Bunches are then generated at the positions determined by the corrugation geometry.

The composite targets, consisting of planar heavy and light ion layers, have also been considered for generation of high quality ion beams [9, 12] required for various applications as hadron therapy [13] and nuclear fusion [14, 15]. When the corrugation is tailored on the interface between two different layers, rather impulsive Richtmyer-Meshkov [16, 17] like instability (RMI) can develop.

These two instabilities belong to the same family group and are being thoroughly investigated [18-21], as they play important roles in several fields like astrophysics (e.g., development of the filament structure of the Crab Nebula [22]), and are affecting the creation of the hot spot in the inertial fusion [23]. The main differences between them are the dependence of the instability appearance on the duration of the driving force and the direction of the acceleration toward the interface [18]. The driving force of RTI is in principle continuous, while RMI is impulsive. RMI can occur when the acceleration is directed toward either side of the interface, whereas RTI can occur only for the direction from lighter to heavier media (LH) [18]. Moreover, in the case of heavy-light (HL) direction of acceleration, RMI exhibits characteristic phase inversion of the corrugated interface as was shown in theory [19, 21, 25] and also in experiments [26]. Several theoretical models of RMI were described including, exact linear solution [25], asymptotic solution [27, 28] and relativistic solution [29]. Behavior of RMI can be also explained by the description of velocity shear induced at the corrugated interface [20, 21].

In this paper we present an effect of a controlled development of collisionless relativistic instability with RMI-like features originating from the interaction of steep-front high-intensity and high-power laser pulse (over $10^{23}$ W/cm$^2$ and 80 PW) with double-layer target with interface modulation. Initially small perturbations at the interface are amplified during the laser-target interaction, leading to the formation of low-density regions at the positions determined by the initial perturbation geometry and high-density plasma bunches between them. The bunches, with higher density than the density of the initial foil are then accelerated by the laser radiation pressure as whole compact structures. Moreover, laser field propagating through the low-density regions creates an enfolding field around the central plasma bunch, prevent-
ing from the perpendicular particle expansion. These behaviors result in the generation of quasi-monoenergetic, well-collimated ion beams with the average energy in the multi-GeV range and transverse emittance of one order of magnitude lower than that in the case of conventional accelerators. The laser accelerated high-energy ion beams from composite targets may also find applications in material sciences and nuclear physics research [30].

The results section of this paper is organized as follows. In the first subsection the mechanisms of the development of the instability with RMI-like features and beam generation is described. In the second one the properties of the collimated quasi-monoenergetic proton beam such as its low emittance and divergence are discussed. In the last subsection our acceleration scheme is compared with schemes with different configurations such as using double-layer target without corrugation, single-layer target, and a case where full-front laser pulse is used instead of the steep-front pulse.

II. RESULTS

To demonstrate the ion acceleration scheme based on double-layer target with corrugated interface between two ion species and its advantages compared to targets without corrugations we perform 2D particle-in-cell simulations using the code EPOCH [31]. The QED (quantum electrodynamics) module [32] resolving non-linear Compton scattering is included in the simulations, since several phenomena from the field of quantum electrodynamics occur in the assumed laser intensity range.

In our case, linearly s-polarized (electric field is perpendicular to the plane of incidence) Gaussian laser pulse with a steep front incidents normally on the target. Its wavelength is \( \lambda = 1 \, \mu m \) and peak intensity is \( I_{max} = 1.37 \times 10^{23} \, \text{W/cm}^2 \), thus yielding dimensionless amplitude \( a_0 = eE_0/m_e\omega \approx 0.85\sqrt{\frac{1}{10^{18}}} \, \text{W cm}^2 / \lambda^2 \, \mu \text{m} \approx 315 \) and critical density \( n_c = e_0m_e\omega^2/\varepsilon \approx 1.15\times 10^{21} \, \text{cm}^{-3} \), where \( E_0 \) is the electric field amplitude, \( \varepsilon_0 \) permittivity of vacuum, \( \omega \) is laser angular frequency, \( m_e \) and \( e \) are electron mass and charge, respectively, and \( e \) is speed of light in vacuum. The amplitude \( a_0 \) can locally reaches values of 500s in our simulations, due to the self-focusing. Therefore, implementing of radiation friction and QED effects [33] is necessary in our case. The laser beam width at the full width at half maximum (FWHM) is 10 \( \lambda \) and beam duration at FWHM equals to 8 laser periods \( T \). The steep front is realized by filtering out the low-intensity part at the front of the laser pulse till 2.4 \( T \) (i.e. 30% of FWHM) before the peak of the temporal Gaussian profile.

This laser profile can be produced by several methods, e.g., using a thin overdense foil, so-called plasma-shutter [34][35] or nonlinear evolution of the laser pulse propagating through an underdense plasma [36][37]. This approach may also improve spatio-temporal contrast of intense laser beam in possible future experiments by filtering a prepulse that accompanies the main pulse [37].

The generated steep front of the laser pulse can reduce the development of transverse short-wavelength instabilities (hereinafter referred as SWI), as proposed in theory [4]. These instabilities cause the disruption of even initially planar foils during radiation pressure acceleration and are usually ascribed to RTI [1][2][3][4] or electron-ion coupled instability [44][45]. Reduction of SWI then enables the development of long-wavelength instabilities induced by the target geometry.

The double-layer targets consists of light and heavy ion layers. The light layer is made of solid hydrogen with electron density \( n_e = 5.36 \times 10^{22} \, \text{cm}^{-3} \), i.e., 48 \( n_c \). It corresponds to targets already demonstrated in experiments [40] (with thickness down to 20 \( \mu m \)). The heavy layer consists of corresponding cryogenic deuterium with the same electron and ion number density, but with two times heavier ion mass than in the case of light layer. Therefore, the Atwood number \( A = (m_2\rho_2 - m_1\rho_1)/(m_2\rho_2 + m_1\rho_1) = \pm 0.33 \), where \( m_{1,2}, \rho_{1,2} \) are the ion masses and densities at the front (1) and rear (2) layers. The sign depends on the direction of the acceleration, i.e., plus for the LH case and minus for the HL case.

The wavelength of the initial sinusoidal interface perturbation is 5 \( \lambda \) and its amplitude is set to 0.25 \( \lambda \). The phase of the modulation is shifted by \( \pi \) between HL and LH cases, assuring the maximum number of proton particles around the y axis. Target thickness is set to 2 \( \lambda \) (1 \( \lambda \) per each layer), corresponding to the optimal thickness \( l \) for radiation pressure acceleration mechanism

\[
\frac{l}{\lambda} = \frac{a_0n_e}{\pi n_c}.
\]

The radiation pressure acceleration mechanism starts dominating over more traditional target normal sheath acceleration mechanism for these relatively low-density (but still overdense) targets, like cryogenic hydrogen, at even lower intensities [42][47][48].

II.1. Mechanisms of the beam generation

Despite its relatively low density, the target does not become fully relativistically transparent, as would be suggested by linear analysis \( n_e < \gamma n_{ec} \) [49] (where \( \gamma = \sqrt{1 + a_0^2/2} \) for linear polarisation), since electrons from the front layer are being pushed into the target by the ponderomotive force, increasing the initial electron density. Therefore, the radiation pressure can still be efficiently acting on the target [47], driving a collisionless compression wave propagating toward the corrugated interface in the middle of the target (see Fig. 1b and 1c), as proposed in theory [13]. For the sake of brevity and continuity with the currently established fluid RMI theory [13], this shock-like jump discontinuity will be hereinafter referred as shock and its reflection in the HL case as rarefaction [19].

In the case of HL interaction, the shock reaches the interface at time \( t = 1.25 \, T \) (see Fig. 1b). Therefore, the average shock speed is above 0.6 \( c \). The incident shock interacts with the corrugated interface (Fig. 1b). The
reflected rarefaction wave is observed in Fig. 1c. Time evolution of phase inversion of the corrugated interface (i.e., switching the positions of corrugation maxima and minima), characteristic for RMI in the HL case can be observed comparing the areas at the left hand side of Fig. 1. Particularly, at the $y$-positions $\pm 2.5 \lambda$ (inversion from corrugation minima in Fig. 1a to maxima in Fig. 1d) and at the $y$-positions $0 \lambda$ and $\pm 5 \lambda$ (inversion from corrugation maxima to minima). The phase inversion results in the stretching of the proton layer as regions of initial corrugation minima stay behind the regions of initial maxima (Fig. 1d). It subsequently creates low density regions between them at the positions where the initial amplitude of the corrugation was zero ($\pm 1.25$ and $\pm 3.75 \lambda$).

Different situation occurs in the case of LH interaction. After the shock hits the interface, the proton layer enters into the deuterium one and eventually propagates through it, as can be seen in the time evolution at the right hand side of Fig. 1. Now the phase on the remaining interface (at $x$-position around $2 \lambda$ in Fig. 1g and -h ) is kept (maxima stay at the $y$-positions $\pm 2.5 \lambda$ and minimum at $0 \lambda$) as predicted by the RMI theory. Deuterium layer becomes (relativistically) transparent to the incident laser pulse and the radiation pressure is acting on the detached proton layer around $x$-positions $3 \lambda$ (Fig. 1g) and $5 \lambda$ (Fig. 1h). The detached proton layer then undergoes phase inversion. However, the driving mechanism is essentially different from that in the HL case and can be explained as follows. The momenta delivered to the particles by the laser pulse in both rectangles presented in Fig. 1c is approximately the same (neglecting effects of Gaussian shape of the pulse around axis). However, the number of protons in the central rectangle (around $y$-position $0 \lambda$) and lateral rectangle differs, i.e., less particles receive the same amount of momenta from the laser pulse in the case of lateral rectangle. Therefore, they can reach higher energies, propagate with higher velocity and eventually overtake the particles initially located inside the central rectangle.

To highlight the positions of the high-energy particles at later time ($t = 47 T$), the energies of the particles are being displayed as colors in Fig. 2 instead of density as in Fig. 1. The laser pulse (depicted with grey color) is also included in this figure. In the HL case (Fig. 2a), the protons on the regions around initial maxima ($0 \lambda$ and $\pm 5 \lambda$) are accelerated to higher energies than around initial minima ($\pm 2.5 \lambda$), which stay behind. Central and lateral bunches are being developed at the positions of initial maxima. The laser pulse propagates through the (relativistically) transparent low-density regions. The propagating electric field creates areas of high ponderomotive potential. Electrons are then pushed to the area of lower ponderomotive potential (i.e., lower electric field $E$) around the axis by the ponderomotive force, subsequently reducing the perpendicular movement of the ions. This field is not enfolding the lateral bunches (around $y$-positions $\pm 5 \lambda$) which then rapidly dissipate in time and only central bunch is present at time $t = 47 T$ as will be shown in Fig. 4a.

In the LH case (Fig. 2b) the central bunch reaches lower energies than lateral areas (around $y$-positions $\pm 2.5 \lambda$), corresponding to the previous discussion about rectangles in Fig. 1e. This situation is opposite to the HL case. The detached proton layer can confine the laser pulse and can be accelerated as bubble [1], till the low-density regions appears in the lateral areas (initial maxima at the $y$-positions $\pm 2.5 \lambda$). Therefore, the maximal reached energy is still slightly higher at time $t = 47 T$ in the LH case (light yellow color in Fig. 2b) than in the HL case (Fig. 2a). However, due to the occurrence of the low-density regions at the positions of initial zeros of the corrugation (HL case) instead of maxima (LH case), more narrow bunch with higher density and also more narrow enfolding field develop in the HL case. This behavior is crucial for long term acceleration as will be shown later.

II.2. Beam energy and Quality

In the HL case, well collimated quasi-monoenergetic proton bunch is developed as can be seen in Fig. 3a, where time evolution of proton energies is shown. Although the energy spread is increasing after time $t = 47 T$, the bunch structure is still kept till the end of the simulation. Moreover, the average bunch energy is gradually shifting closer to the maximum energy during time. The average beam energy at time $t = 47 T$ reaches 1882 MeV and bandwidth (at FWHM) is 69 MeV (Fig. 3b). Therefore the energy spread is about 3.7 %. Angular distribution of the ions in the beam (red part of Fig. 3b) is shown in Fig. 3c. This graph implies low angular spread of $2\theta = 0.65^\circ$ (at FWHM). Therefore, the solid angle is $\Omega = 2\pi (1 - \cos(\theta)) = 0.1$ mrad. The normalized rms transverse emittance of these particles is $\epsilon_{rms} = [\epsilon^2 + (p_x/p_y)^2]^{1/2} = 0.046 \text{ mm mrad}$.

Where $m_p$ and $p_y$ are proton mass and momentum in the (transverse) $y$-direction. $\epsilon_{rms}$ is proportional to the area of the ellipse containing particles in the phase space ($y$-$p_y$). Referring definition and further discussion can be found in Ref. [50]. This emittance is one order of magnitude lower than in the case of conventional proton accelerators [51, 52], but still one order of magnitude higher than the emittance reported in Ref. [51] for much lower energy range of protons up to 10 MeV.

The transverse emittance can be also defined in real space via the beam divergence ($\Theta_{div} = \sqrt{\sum_{i=1}^{N} (\Theta_i - \langle \Theta \rangle)^2/N}$) as in [53]

$$\epsilon_y = \frac{4}{N} \sqrt{\sum_{i=1}^{N} (y_i - \langle y \rangle)^2} \times \sqrt{\sum_{i=1}^{N} (\Theta_i - \langle \Theta \rangle)^2}. \quad (2)$$

That yields the values for protons at the FWHM of the bunch (Fig 3b): $\Theta_{div} = 0.038$ rad and $\epsilon_y = 0.218 \text{ mm mrad}$. Assuming the whole high energy
proton beam (from 1748 MeV, i.e., the beginning of the bunch waist, to the maximum beam energy of 2770 MeV) the values rise to \( \Theta_{\text{div}} = 0.051 \) and \( \epsilon_u = 0.372 \text{mm} \cdot \text{mrad} \). This transverse emittance is still two order of magnitude lower compared to the one in Ref. [53], in which protons reached similar energies to our case (1.67 GeV) in the simulation with underdense hydrogen target.

Approximately \( 0.32 \times 10^6 \) simulation particles have energy within the assumed energy range, which means about 6.7\% of \( 4.8 \times 10^6 \) simulation particles initially located inside the laser spot area (i.e. \( \pm 5 \lambda \)). That yields \( 1.43 \times 10^{12} \) real particles and charge of 229 nC, assuming cuboid target with the size of \( 40 \mu \text{m} \) in the \( z \) direction. Around 40.4\% of the laser pulse energy is converted into the particles propagating in the forward direction (protons: 28.7\%, deuterons: 6.1\%, electrons: 5.6\%). Around 3.4\% of the laser pulse energy is converted into the proton particles in the energy range of 1882 ± 69 MeV (corresponding to the FWHM red part of Fig. 4b).

II.3. Comparison with similar types of targets

To point out the effects of two layer target and modulation on its interface, another two simulations are performed. Namely, two layer deuterium-hydrogen target without corrugation, referred as HL-WO, and pure hydrogen target of the same thickness (2\( \lambda \)), referred as L2. The reference time was chosen to be \( t = 47 \text{T} \) as in the previous discussion. Proton energy spectra of all the simulated cases are shown in Fig. 3d and corresponding proton density and energy spatial distributions are presented in Fig. 4. In the HL and LH cases (with the interface modulation) a bunch structure is created in the spatial plot (Fig. 4a and -b) and corresponding quasi-monoenergetic peak is developed in the proton energy spectra (Fig. 3d). The trend of a more narrow bunch in the HL case compared with the LH case, initialized by different positions of the low-density regions of the foil, also continues in the later time. This results into more narrow peak in the energy spectra with higher average energy in the HL case (1890 MeV compared to 1490 MeV) in Fig. 3d. Maximal reached energy is also slightly higher in the HL case (2770 MeV compared to 2700 MeV). Proton bunch is more shifted in space towards the front of the overall proton cloud and less spread in the\( x \)-direction in the HL case (compare Figs 4a and 4b).

On the contrary, HL-WO and L2 cases develop a bubble structure (Fig. 4c and 4d), with no significant peaks in the energy spectrum (Fig. 3d). The maximum proton energy of these simulated cases is reached in the L2 case (2850 MeV). This case is the best match for the RPA regime described by Eq. 1. Assuming single layer target with no induced instability. The symmetric bubble structure holds till the reference time and proportionally small amount of protons near the axis are accelerated to the highest maximal energies (Fig. 4d). Oppositely, the bubble shape is getting distorted in the HL-WO case (Fig. 4c), due to the multi-ion-species effects, resulting into significantly lower maximum energy (2430 MeV).

At the same reference time deuterium ion reached significantly lower maximal energies than protons in our simulations (see inset in Fig. 3d). Their energies are lower than 1200 MeV (i.e., 600 MeV per nucleon) and the highest energy is reached in the LH case. It corresponds to the observation of proton bunch being detached from the deuterium layer, reducing further acceleration of deuterium ions.

III. DISCUSSION

In conclusion, collisionless relativistic instability with RMI-like features is observed in our simulations, using double-layer targets with initial interface modulation and high-intensity steep-front laser pulse. Evolution of this instability is described in the heavy-light (HL) and light-heavy (LH) cases, resulting into a development of bunch structures in the density distribution and proton spectra.

Well-collimated, quasi-monoenergetic proton bunch is observed for the HL case, with the average energy in the multi-GeV range, energy spread down to 3.7\%, solid angle of 0.1 mrad, divergence of 0.038 rad and transverse
rms emittance down to 0.046 mm·mrad (at FWHM). The emittance is one order of magnitude lower compared to that of conventional accelerators. The energy conversion of the laser pulse into the proton particles inside this bunch is around 3.4 % and overall conversion into all particles is over 40 %.

The bunch is more distinctive in the proton spectrum and its average energy is significantly higher compared to other simulated cases with the same target composition. These other cases are either dominated by essentially different type of instabilities (development of short-wavelength instability using full-front laser pulse or introduction of the modulation on the surface of a single layer target) or the instability and bunch structures are not developed (targets without initial modulation using steep-front laser pulse, reducing the development of short-wavelength instability).

IV. METHODS

IV.1. Simulation setup

The laser plasma interactions occur in the simulation box with the size of $80 \times 40 \lambda$. The mesh has square cells. The size of the cells is set to 0.01 $\lambda$ to be shorter than the plasma skin depth $c/\omega_{pe} \approx 0.02 \lambda$, where $\omega_{pe}$ is electron plasma frequency. Since the 3rd order b-spline shape of the simulation particles and current smoothing are used in our simulations, it is ensured that numerical heating is strongly reduced even for the cells larger than the plasma Debye length $\lambda_D$. The simulation time step has been set by EPOCH code in order to satisfy CFL (Courant-Friedrichs-Lewy) condition \cite{54} to $6.7 \times 10^{-3} \ T$. Each cell inside the plasma slab initially contains 48 simulation electrons and the same number of protons or deuterium ions, respectively. Temperatures of all particles are initialized to 5 keV to further reduce numerical heating and the particle solver begins to move the particles just a few time steps before the arrival of the laser pulse front to the target. Target is placed at the position $x = 0$, situated $10 \lambda$ from the simulation box boundary in the direction of laser propagation. The time instant, when the laser pulse front reaches the edge of the plasma is referred as $t = 0$. The transverse size of the target is $40 \lambda$, i.e., the target is touching simulation box boundaries at positions $y = \pm 20 \lambda$ where thermal boundary conditions for particles are applied.

IV.2. Data visualisation and virtual reality

Several visualization methods were used to create visual outputs from the above discussed simulation data. Firstly, raw data were imported in ParaView \cite{55} software to interpret them and to create visual outputs (figures and videos) describing the simulation mechanisms. Secondly, data have been presented in web-based interactive 3D application (see Fig. 6) which runs in a regular web browser and utilizes VR mode to offer scientists completely new point of view of their simulations. This second option comes together with a challenge to find the best workflow that would enable visualization of such a large datasets in a web browser while maintaining high frame rates for smooth experience not only while using the VR headset. We have evaluated existing technologies like VTK.js library developed directly for scientific visualization in browser as well as Three.js javascript 3D library. Unfortunately, both libraries could not render our datasets with hundreds of timeframes with sufficient performance. The main issues were minutes-long loading times together with insufficient frame rates for large animated datasets. Thus, we have developed custom WebGL solution, a framework that not only renders the dataset on the GPU in real-time at high frame rates, but also provides orthogonal views, textual and numeric information, alongside a GUI containing timeline animation controls and layer visibility management, with additional graphical elements based on D3.js for plotting animated graphs and legends. In order to import simulation data to this application, they are transformed into binary buffers in a node.js script so the visualization engine can directly send them to GPU with no need for further pre-processing. Currently we are working on an implementation of our transformation pipeline to the Jupyter notebook \cite{56} toolchain to allow more users to create high performance web based VR enabled visualizations.

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VI. AUTHOR CONTRIBUTIONS

M. Matys carried out the simulations, analyzed the results, generated the graphs and wrote the bulk of the manuscript. M. Kecova generated the figures and supplementary video and prepared the part of the methods section about the virtual reality. K. Nishihara, J. Psikal and S. V. Bulanov provided theoretical interpretation of the results. K. Nishihara also contributed for determining simulation parameters and analysis of the instability with RMI-like features. S. V. Bulanov with G. Korn also provided overall supervision. All authors contributed to the preparation of the manuscript.

VII. COMPETING INTERESTS

The authors declare no competing interests.
FIG. 1. The instability development in the cases of HL (left column) and LH (right column). Blue and red scales represent deuterium and proton densities with maximum value set to the initial density. Full density is indicated by the vertical height. Time instants at rows: 1.25 T, 3 T, 5.5 T and 8 T.

FIG. 2. The ion local mean energy, represented by red and blue color scales, their density, represented by the vertical height and the laser pulse electric field in the z-direction (represented by both the vertical height and by the grey scale) at time $t = 14 T$. Simulated cases: a) HL, b) LH.
FIG. 3. a) Time evolution of the tail of the proton energy spectra in the HL case target, b) proton energy spectra of HL target at time $t = 47 T$, with highlighted FWHM section used for c) angular distribution, d) proton energy spectra (corresponding deuteron energy spectra in inset) for various targets (see details in the text) at time $t = 47 T$. 

FIG. 4. The proton density (represented by the vertical height), proton energy (represented by the blue to yellow scale) at time $t = 47 T$. The simulated cases are: a) HL, b) LH, c) HL-WO, d) L2 (see parameters of all the targets in the text). 

FIG. 5. a) Short-wavelength instability in the HL-FF case at time $t = 14 T$. Blue and red scales represent deuterium and proton densities with maximum value set to the initial density. Full density is indicated by the vertical height. b) Proton and deuteron energy spectra for HL, HL-FF and L2-SM cases at time $t = 47 T$. 
FIG. 6. Data visualisation in web-based interactive 3D application.