Relativistic laser nano-plasmonics for effective fast particle production

A Andreev, K Platonov, J Braenzel, A Lübecke, S Das, H Messaoudi, R Grunwald, C Gray, E McGlynn and M Schnürer

1 Max-Born Institute, Berlin, Germany
2 ELI-ALPS, Szeged, Hungary
3 Sankt Petersburg State University, St. Petersburg, Russia
4 State Technical University, St. Petersburg, Russia
5 School of Phys. Sci., National Centre for Plasma Science and Technology, Dublin, Ireland
6 Technical University Berlin, Straße des 17. Juni 135, 10623 Berlin, Germany

Abstract
We have studied particle acceleration in different nanostructured targets irradiated by high intensity laser pulses of high contrast. We find that the maximum energy of emitted particles and their directionality is significantly enhanced in the case of nanostructured targets with respect to plane targets. We have studied theoretically in detail the generation and propagation of fast electrons in nanowire targets. Such targets exhibit an extraordinary high conversion efficiency of laser energy into electron kinetic energy. We observe guiding of electron bunches along the wires. Results from theory and simulation compare reasonably well with the experimental data.

Keywords: high intensity laser, laser plasma, nanostructure targets

Introduction
Laser driven nano-plasmonics deals with optical processes in plasmas at relatively low intensities and on nanoscale, i.e. on the order of or smaller than the wavelength of the laser radiation [1]. The underlying physical process is connected with a laser field enhancement due to proper nano-structuring of a target. Laser-matter interaction including nanoscale confinement of radiation and its conversion provides attractive opportunities for both, fundamental research and technological applications [2].

Laser nano-plasmons produced by structured surfaces [3] have recently gained growing attention in laser plasma particle acceleration physics, as it can significantly improve the acceleration mechanism. Today’s research on laser driven particle acceleration focuses on the development of efficient and compact particle sources for different applications. Basically, a laser at relativistic intensity interacts with thin foils [4] generating fast electrons, which subsequently accelerates ions of low emittance up to high energies [5]. Nevertheless, up to now the acceleration efficiency is quite low, even if the target thickness is optimized with respect to the laser parameters. The development of new targets pushing the conversion efficiency is the central goal of current research activities [6].

It is obvious, that surface modification acts back on the light absorption in general. Therefore, surface nanostructuring is a promising strategy for developing a new target design that allows for higher conversion efficiency in particle acceleration [7]. A key question concerns the persistence of the structures during the high intensity laser irradiation and its consequences. Apparently, the laser properties and optimum target structure are linked together. In this paper we will study this relation, in detail.

In the following, we investigate the interaction of an ultra high intensity and high temporal contrast (<10^{-16}) laser pulse with different nanostructured targets. Special emphasis is placed on the generation of fast ions. We observe that their maximum energy and directionality can be significantly enhanced by specific nanostructures. Under optimum conditions the absorption of laser energy can approach 100%, and
the conversion of laser energy into ion kinetic energy is significantly enhanced. Our results from theory and simulation are in reasonable agreement with our experimental findings for different nanostructured targets.

**Formation of nanostructures by femtosecond laser**

Recent theory predicts a new class of meta structures by engineered sub-wavelength entities, which enable unprecedented properties for electro-magnetic radiation [1]. These specific structures do not exist in nature and will have profound impact on integrated nanophotonics. One example is the realization of a so-called super lens that focuses below the diffraction limit [2]. Considering practical applications, easy-to-handle and cost-effective methods for surface structuring using different materials are desirable. One of the most convenient top down approaches is the formation of laser-induced periodic surface structures (LIPSS) [8], which are regarded as a self-organizing mechanism. It exploits the interference of the incoming light wave with a secondary wave (scattered light or a surface plasmon wave generated in the material) and its propagation along the surface [9]. At high intensities and broadband excitation as delivered by ultrashort-pulse lasers, the scenario becomes more complex by including multiphoton excitation channels and/or spectral selectivity. As shown for different types of materials like, e.g. W [10], Si [11], TiN [12] or diamond-like carbon [13], nonlinear excitation with femtosecond laser pulses enables structural periods well below the excitation wavelengths (sub-wavelength LIPSS). For the formation of such nanoscale LIPSS (nanoripples) in dielectric media, two distinct types of structures are observed, which can be described as ablation reliefs (see, e.g [12]) or the creation of deep, nearly planar grooves or nanoplanes normal to the surface. Extended-area structuring on fixed and moving substrates was obtained by applying line focusing with a cylindrical lens. Under defined conditions with respect to pulse number, pulse energy and scanning velocity, two types of ripple-like LIPSS with high and low spatial frequencies corresponding to periods in the range of 90 nm and 340 nm, respectively, were formed [13]. In the present work, we experimentally investigated the effect of LIPSS nanostructures on a titanium surface on laser ion acceleration. LIPSS structures were generated as described in [13]. Besides those LIPSS, we also investigated a second class of nanostructures with significantly smaller structural width, i.e. nanowires. These ZnO nanorods were grown on a Ti foil by chemical bath deposition on drop coated seed layers [14, 21].

**Interaction of intense short laser pulse with nanostructured targets**

**Absorption**

Both, experiment and theory show a significant enhancement of the maximum ion kinetic energy and the ion flux in laser driven ion acceleration if very thin plane target films rather than thicker samples are used. Nevertheless, the absorption of high contrast laser pulse in the range of $10^{18} \leq \rho \leq 10^{22} \text{ W cm}^{-2}$ does not exceed 50% in thin flat targets (<100 nm) [6]. An ideal target combines maximum absorption and minimum mass which can be realized by a thin substrate foil covered by long nanowires, i.e. so-called ‘brush’ or ‘snow’ targets [16, 23]. The absorption of laser energy in such a target reaches up to 100% and the conversion of laser energy into ion kinetic energy can approach tens of percent [17]. In the following, we analyze the efficiency of the interaction of grating-like nanostructures on a foil target with an ultraintense laser pulse.

In our simulations we investigate a substrate foil of few hundred nanometer thickness at solid density. An imposed relief structure on front and rear side is characterized with different ledges. The parameters of the rectangular ledge at the front side are denoted as: $d_1$—width, $d_2$—distance between the ledges, $h$—ledge height and the thickness of the substrate foil as $s$, see figure 1. The parameters of the relief at the rear side are denoted just as those at the front side but with capital letters.

We also varied the density of ion constituents, i.e. we considered heavier ions (C$^+$) and light ions (H$^+$). In addition, protons can be deposited as a thin layer on the rear surface. The purpose of target parameter optimization is to achieve both, a high absorption of the laser pulse in the target and the maximum energy of the accelerated ions.

**Simulation model**

The investigation is based on 2D particle-in-cell (PIC) modeling of the motion of plasma particles. With the help of the
modified code PSC, the modelling was performed for different targets. In 2D simulations we used the following parameters: a laser wavelength of $\lambda_L = 0.8 \, \mu m$, intensity values of $I_L = (0.5, 1, 5) \times 10^{20} \, W \, cm^{-2}$, pulse durations of $t_L = 15, 30 \, fs$ and a spot size $d_L = 3 \, \mu m$. The simulations have been performed in a simulation box of $100 \times 100 \, \mu m^2$ divided into $5 \times 10^7$ cells with 30 particles of each type in a cell.

It was shown [18] that in case of a target with a non-structured rear side the electric field at the rear side has a regular profile. Contrary, at the structured front side, the field reproduces a relief distribution. Thus, the quality of the accelerated proton beam differs for each side. We will now consider only the rear side component.

In figure 2(a) black dots represent the simulation results on the absorption enhancement of the nanostructured target relative to a plane foil as a function of parameter $d_1$. For an optimum structure, the absorption is enhanced by a factor of 6 at a laser intensity of $10^{20} \, W/cm^2$. Absorption of the flat target was only 14%, thus the relief enhances the absorption to 84% at the chosen laser pulse and target parameters. It is worth to note that plasma gradient can enhance absorption as well, but this is a function of the laser intensity at $t_L = 30 \, fs$. The simulations have been performed in a simulation box of $100 \times 100 \, \mu m^2$ divided into $5 \times 10^7$ cells with 30 particles of each type in a cell.

Analytical model

To explain these dependencies we developed an analytical model based on our previous work [19]. According to this model the absorption mechanism is similar to Brunel’s one: an electrons are extracted from target surface and fill relief dips in a given laser cycle. During the next laser cycle they are pushed back into the target. Those electrons absorb the main part of laser energy and transport it inside target. Therefore, the number and the integral energy of these electrons are calculated and the absorption coefficient is estimated from the ratio of the integral electron energy to laser energy. The laser field acting on the surface of the relief with an orientation perpendicular to the laser polarization is capable of taking electrons into vacuum from the maximal extraction depth of $\xi_{\text{ext}} = E_L/(4\pi l_e)$, where $n_e$ is the electron density in the skin layer and $e$ is the electron charge. Therefore, $d_L \geq 2\xi_{\text{ext}}$ should be fulfilled. The number of electrons leaving the unit area of the relief at optimum conditions is $N/e = n_d/d_L^2 = \frac{(d_L/2\xi_{\text{ext}})^2}{1 + (d_L/2\xi_{\text{ext}})^2}E_L/e$. The total lateral surface, from which electrons leave, is $S = h_2L/(d_1 + d_2)$, where $h_2$ is the laser spot-size. The characteristic size of an electron orbit near the relief is $\rho_{\text{eh}} = C e E_h L/m_e\omega_L^2$, where constant $C \leq 1$ (for a single electron in vacuum $C = 1$, but the influence of other electron decreases $C$). Therefore, the optimum distance between the two ledges of the relief can be written as, $d_2 = 2\rho_{\text{eh}}$. Accordingly, the maximum electron energy in vacuum can be estimated for an optimum profile by:

$$\varepsilon_{\text{tot}} = \sqrt{\frac{d_L}{2\xi_{\text{ext}}}} \frac{(d_L/2\xi_{\text{ext}})^2}{1 + (d_L/2\xi_{\text{ext}})^2}(h_2E_h L/(d_1 + d_2))\rho_{\text{eh}},$$

where $m_e$ is the electron mass, $\omega_L$ is the laser frequency, $\varepsilon_{\text{eh}} = m_e\rho_{\text{eh}}^2/2$. Let’s mark that $\varepsilon_{\text{tot}}$ is the additional energy absorbed by relief. For flat target ($h = 0$) full electron energy is determined as $\varepsilon_{\text{tot}}(\text{pl}) = \eta\varepsilon_{\text{eh}} = \eta (E_h^2/4\pi L(d_1/2\xi_{\text{ext}}))^2(h_2),$ where $\eta$ is absorption coefficient of flat target. Then, the relative absorption is given by:

$$\chi = \frac{\varepsilon_{\text{tot}}(\text{pl}) + \varepsilon_{\text{tot}}(\text{nl})}{\varepsilon_{\text{tot}}(\text{pl})} \frac{1 + (d_L/2\xi_{\text{ext}})^2}{\varepsilon_{\text{tot}}(\text{pl})} = \begin{cases} \frac{1}{\chi_{\text{max}}}, & \chi > \chi_{\text{max}} \\
\frac{2\varepsilon_{\text{tot}}(\text{pl})}{(d_1 + d_2)} + \frac{4\varepsilon_{\text{tot}}(\text{pl})}{\sqrt{2}\varepsilon_{\text{tot}}(\text{pl})}, & \chi \leq \chi_{\text{max}} \end{cases}
$$

(1)

Here $\chi_{\text{max}}$ is the relative absorption for the optimum relief when the laser energy $\varepsilon_L \approx \varepsilon_{\text{tot}}$ which corresponds to simulations (see figure 2(a)). With (1) the dependency $\chi$ on parameter $d_1$ is calculated and plotted with a red solid line in figure 2(a). A good agreement with simulation results has been obtained for $C = 0.6$ at $\eta = 0.14$, taken from the simulation.
In the following we consider the influence of the rear side relief on the parameters of accelerated particles. It was shown [18] that only at the front side of the target the ambipolar electric field reproduces a relief distribution. Thus, hot electrons have a quasi-equilibrium state in the ion field and can be described with the help of a density distribution at an effective temperature because of electron re-circulations in a thin foil. Therefore, the influence of the front relief is described by varying the temperature and density of hot electrons. Laser pulse parameters are taken into account via expansion of a relativistic electron gas inside the target and its relevant parameters. During acceleration, the laser irradiated target expands, and the parameters of the electron gas obey the adiabatic law, \( p_x n_e^2 = T_0 n_0^2 \), where \( p_x \) and \( n_e \) are the pressure and density of hot electrons, respectively, and \( T_0 \) and \( n_0 \) are the initial values of their temperatures and density. Since the acceleration time of ions is longer than the traveling time of a relativistic electron through the target, we can neglect the inertia of electrons in the hydrodynamic equation of motion. We write the condition of local equilibrium of the electron gas as \( \phi \partial^2 \phi / \partial r^2 = -n_e \partial p_e / \partial r \), where \( \phi \) is the scalar potential of the ambipolar field in the target. Using the adiabatic law of expansion of the electron gas, one can easily express the electron density from the balance equation of motion in terms of the potential as follows: \( n_e(\phi) = n_0(1 + \frac{\gamma - 1}{\gamma} \frac{\phi}{T_0})^{-\frac{1}{\gamma - 1}} \). We choose the adiabatic index as \( \gamma = 2 \) because of 2D geometry in our simulations. Therefore, the 2D Poisson equation for the potential and the field that accelerates ions becomes linear \( \Delta \phi = -4\pi \frac{|\mathbf{e}|^2}{c^2} |Zn_z(y, z) + n_{0z} - n_e(\phi)| \) and can be solved for the relief target. The density of the positive heavy ions for the rear relief (see figure 1) can be written as \( n_{1z}(y, z) = n_{10} \sum_k \theta(y - k(D_2 + D_1)) \theta(kD_1 + (k + 1)D_1 - y) \cdot \theta(H/2 - |z|) + n_{10} \theta(-H/2 - z) \theta(z + H/2 + s) \) where \( \theta(z), \theta(y) \) are step functions. The density of light ions \( n_{20} < < n_{10} \) is described as an ultrathin layer at the surface of relief (contamination thickness \( \epsilon < D_1 \)): \( n_{2z}(y, z) = n_{20} \sum_{k=\infty}^{\infty} \delta(y - k(D_2 + D_1)) \times \delta(kD_1 + (k + 1)D_1 - y) \cdot \theta(H/2 - |z|) + \theta(y - k(D_2 + D_1)) \theta(kD_2 + (k + 1)D_1 - y) \cdot \delta(H/2 - z) + \theta(k(D_2 + D_1) - y) \cdot \delta(y - (k - 1)D_1 - kD_1) \delta(-H/2 - z) \). Let us introduce the dimensionless potential \( \psi = |e| \frac{\phi}{T_0} \) and normalize the coordinates \( \xi = y/d_0 \) (along the target surface), \( \zeta = z/h \) (perpendicular to the target surface) and the time to the ion plasma frequency \( \tau = \omega_p/\omega_{pe} = 4\pi Z^2 n_0 e^2/m_e \). In dimensionless variables, the Poisson equation for the dimensionless density of hot electrons, which scales linearly with the potential \( \psi = 2(\eta_e - 1) \), takes the following form: \[
 2\left( p_e \frac{\partial^2 \phi}{\partial \xi^2} + \rho \frac{\partial^2 \phi}{\partial \zeta^2} \right) = \eta_0 n_0 - \sum_k \theta(\xi - k(1 + D_1/D_2)) \times \theta(k + (k + 1)D_1/D_2 - \xi) \cdot \theta(1/2 - |\zeta|) + \theta(-1/2 - \zeta) \times \theta(\zeta + 1/2 + s/H) \]
(2)
The equation of proton motion can be written as:
\[
 \frac{\partial^2 \xi}{\partial \tau^2} = \frac{2p_e^2 \partial n_e}{\eta_0} - \frac{2p_e^2 \partial n_e}{\eta_0} = \frac{2p_e^2 \partial n_e}{\eta_0} \frac{\partial^2 \phi}{\partial \xi^2} = \frac{2p_e^2 \partial n_e}{\eta_0} \frac{\partial^2 \phi}{\partial \xi^2} = \frac{2p_e^2 \partial n_e}{\eta_0} \frac{\partial^2 \phi}{\partial \xi^2}.
(3)
\]
Here \( r_0^2 = T_0 4\pi e^2 Z n_0 \), \( \eta_0 = n_0 / Z n_0 \), \( \rho = r_0 / d_2 \), \( \mu = r_0 / h \).

Since \( n_{20} < < n_{10} \) is valid protons are considered as test particles moving in the field of heavy ions and hot electrons. The linear differential equation (2) was solved with help of a 2D Fourier transformation.

The potential inside the relief forms a saddle with a convex structure along the \( \zeta \)-axis and a weakly concave structure along the \( \zeta \)-direction. Owing to this saddle shape, the light ions starting inside the relief are squeezed towards the central axis at \( \zeta = 0.5 \) (see figure 3(a)) slide down the potential surface and being finally ejected out of the structure. The numerical solution of the equation of ion motion (equation (3)) in the potential (equation (2)) is shown in figure 3(a).

The proton trajectory is determined by its initial position \( (\xi_0, \zeta_0) \) in the thin surface layer (figure 3(a)). The number (per unit length) of protons accelerated from the lateral relief surface in figure 3(a) and located in the interval \( d_0 \) is \( dN = n_{20} d_2 \xi d_0 \), \( \xi_0 \in [0; 1] \), \( \zeta_0 = 0 \). The number of protons accelerated from the horizontal surface located in the interval \( d_0 \) is \( dN = n_{20} d_2 \xi d_0 \), \( \xi_0 \in [1 + D_1/D_2], \zeta_0 = 0 \).

Dividing this number by the angular interval \( d\theta \) is the proton emission angle, see figure 3(a), one can get \( dN/d\theta \), thus the parametric dependence (parameters \( \xi_0, \zeta_0 \)) of the proton angular distribution function is as follows:
\[
 dN/d\theta = n_{20} d_2 H d_0 d_0 \frac{d\theta(\xi_0)}{d_0} + n_{20} d_2 D_2 \frac{d\theta(\zeta_0)}{d_0}.
\]
From (4) it is obvious that the maximum of the angular distribution is at \( d\theta(\xi_0) / d_0 \approx 0 \), \( \zeta_0 \approx 0 \).5 (see figure 3(a)) and therefore we find its maximum at \( \theta \approx 0 \).5. In figure 3(b) the distribution (4) is represented by a black solid line for the following parameters: \( D_1 = 100 nm, D_2 = 540 nm, H = 200 nm, n_0 = 4 \cdot 10^{20} \text{cm}^{-3} \), \( m_{10} = 6 \cdot 10^{-22} \text{cm}^{-3} \).4 Andreev target was 300 nm. Laser parameters were: \( L_1 = 10^{19} \text{W cm}^{-2} \), duration \( t_1 = 30 \text{fs} \) and the spot size \( d_1 = 3 \text{\mu m} \). Angular and spectral characteristics of accelerated protons obtained with and without a relief are given in figure 4.
Figure 4 shows that the optimum front relief increases the number of protons by a factor of 2 and the maximum energy by a factor of 1.6. The width of the angular distribution ($\pm 10^\circ$) is about the same as for the plane foil. The rear relief slightly increases this width due to the additional side (lateral) surfaces (blue line of figure 4(a)). It is seen that the formula (4) of the analytical model agrees well with the simulation results (see green and blue curves). Targets with front and rear nanostructures and those with structures only at the front side yield very similar maximum proton energies and the same order of magnitude of integrated proton energy (see figure 4(b)). However, the rear relief significantly increases the number of protons at high kinetic energies (4–16 MeV) as it has a larger surface providing an enhanced ambipolar field for proton acceleration. It is worth to note that figure 3 shows a similar effect as in paper [24], but for a rectangular relief profile. While our work focuses on the relief influence on proton numbers and energy distribution, ref. [24] investigates the relief influence on proton spatial distribution and its emittance. Thus, the optimal relief at the foil rear enhances the total number and energy of protons keeping its emittance at the level of plane foil.

Stability of the relief structures

The previous results show that relief structures play an important role in laser absorption and in consecutive particle acceleration. Therefore, it is necessary to investigate whether these relief structures are preserved during the full interaction time with a high intensity laser pulse. We divide the investigation of the relief stability into two parts: At first, we study the interaction in the case that the target does not move in the laser...
propagation direction. And afterwards we will include this motion. For a stationary target, the pre-pulse of the main laser pulse can destroy the surface modulation due to heating and subsequent pre-plasma formation. It is possible to estimate the laser pre-pulse duration, which is not detrimental for the surface modulation, i.e. at which the spatial extent of the ion thermal wave is smaller than the spacing between the ledges. The intensity of the pre-pulse corresponds to the experimentally achievable contrast. Estimation of absorption at a certain contrast level and following target expansion translates into a restriction for the pre-pulse duration. In case of a light target material, typically the duration of nanosecond is derived. Replacement of carbon by heavy material can increase the tolerable pre-pulse duration up to several nanoseconds. So far, the discussion concerned a stationary target. We will now turn to the case that the target can move during laser interaction. The ponderomotive pressure of the laser pulse can deform the relief, if electrons are pushed deeply inside the target. The resulting strong electrostatic field causes deformation of the target surface if the laser intensity approaches $\sim 10^{21} \text{ W cm}^{-2}$ for our target parameters [18].

The question of limitations due to the finite pulse duration is also very important. In figures 5(a) and (b) the spatial distribution of the relief is shown at different times. Apparently, ion expansion from the initial distribution generates a new dynamic structure exhibiting a maximum density at the positions of previous minima. This dynamically formed new relief persists sufficiently long (~100 fs), and efficient absorption of incident laser light at very high intensity is possible even for several 100 fs, which makes this scheme interesting even for longer pulses. In order to approximate the dynamic lifetime of the relief we use the model concerning equations (2) and (3) and insert parameters $d_1, d_2, h$ instead of $D_1, D_2, H$. Smoothing of a relief (increase of $d_1$) can be interpreted as the motion of its vertical (lateral) boundary in the potential $\psi(\xi, \zeta = 0)$ (see figure 3(a)). The motion of relief boundaries obeys the period conservation $d_1(t) + d_2(t) = d_{10} + d_{20} = d_0$. A relief reproduces itself if $d_1 > d_{10} + d_{20}$ resulting in new dense ledges. We associate the relief lifetime $\tau$, with the time when $d_1(t_\tau) = d_{10} + d_{20}$. Considering the motion of potential boundaries of the relief, one can write the following formula for the relief lifetime:

$$\tau = \int_0^{1/2} \frac{\text{d} \xi}{\sqrt{2(\psi(\xi = 0, \zeta = 0) - \psi(\xi, \zeta = 0))}} \quad (5)$$

The calculation (5) for the parameters used to obtain the results in figure 5 yields about 100 fs which is close to 70 fs when in figure 5(b) the collision of ion fronts from side walls of the front relief are observed. The lifetime of this new structure (dynamic relief) is about 100 fs allowing an irradiation with pulses of several 100 fs duration.

**Fast particle generation**

Now, after having studied target dynamics, the maximum proton energy attainable with nano relief targets will be investigated.

In figure 6(a) the dependence of maximum proton energy (red squares) on the laser intensity is shown for optimum relief parameters. For comparison, the maximum proton energy for a plane foil is shown (green squares). The red cycle shows the experimental result for ripple target at $d_1 = 200 \text{ nm}$, $d_2 = 100 \text{ nm}$, $h = 200 \text{ nm}$. Because it was non-optimal relief the proton energy locates between red and green squares. Based on the simulations for an optimum relief we derive scaling equations for the maximum proton energy (above 5 MeV) and the number of light and heavy ions:

$$\varepsilon_p \approx 17(I_l/5 \cdot 10^{19}\text{Wcm}^{-2})^{0.61} \text{[MeV]}$$
$$N_p \approx 7 \cdot 10^{7} \cdot (I_l/5 \cdot 10^{19}\text{Wcm}^{-2})^{0.86}$$

$$\varepsilon_C \approx 69(I_l/5 \cdot 10^{19}\text{Wcm}^{-2})^{0.62} \text{[MeV]}$$
$$N_C \approx 7.4 \cdot 10^{7} \cdot (I_l/5 \cdot 10^{19}\text{Wcm}^{-2})^{0.75}$$

**Figure 5.** (a) The initial profile with a front- (down) and rear-side (up) relief of the target (b) creation of the new dynamic relief when stagnation of ion fronts from side parts of the relief occurs at $t = 69 \text{ fs}$ after interaction of a 30 fs laser pulse. The laser and target parameters are shown in figure 4.
The main part of the absorbed energy is transformed into kinetic energy of the heavy ion component $\varepsilon_C$. This energy component is calculated using energy and momentum conservation laws: in the lab frame one obtains

$$\varepsilon_C - m_Cc^2 = \frac{2\varepsilon_L - F(h_i)}{N_C(h_i)} \frac{(2\varepsilon_L - F(h_i))}{2m_Cc^2N_C(h_i) + 2(2\varepsilon_L - F(h_i))}$$

with $R, \eta$ being the reflection and absorption coefficients, respectively, taking into account that for a nontransparent target $R + \eta = 1$. Incorporating simulation results we yield $N_C(h_i) = I(h_i)I_{\alpha}^{\beta/2}$, $I_L < I_{\alpha}$; $N_C(h_i) = N_i$, $h > I_{\alpha}$ where $N_i$ is the total ion number in the laser spot and $I_{\alpha} = cE_{\text{crit}}^2/(4\pi) = \pi e^2\rho_i^2$ is the critical laser intensity at which all electrons are removed from the target and Coulomb explosion occurs instead of TNSA acceleration, i.e. at which the laser electric field equals the electrostatic field of the ion core $E_L = eN_i$.

The energy loss function describing the energy flux directed outward from the laser spot is:
The depend-
ence \( \varepsilon_c = \varepsilon_c(l_h) \) from (6) is shown in figure 6(b). Equation (6) describes average ion energies and explicitly not the maximum values. These values are quite similar in analytical and numerical calculations. At very high laser intensities \( I_h > 10^{22} \text{ W cm}^{-2} \), in the so called ‘piston’ regime, the absorption coefficient is close to unity and a surface structure does not provide any longer an advantage.

**Transport of hot electron current**

Transmission of high electron currents in targets is always accompanied by generation of cold return currents, which compensate self-generated magnetic fields of the beam. In dense targets such return currents lead to instabilities, filamentation of the relativistic current and turbulences. In targets of specific structures such processes can be mitigated, because the parameters of filamentation are known and the process may be suppressed with an appropriate transport structure. If the propagation length of the relativistic electron beam can be significantly extended, this will also become interesting for potential applications. Present technologies allow engineering targets with a bundle of wires of nanometer diameter [20], which will not be destroyed by the laser pulse with sufficient high temporal contrast [21]. The motion of the hot electrons along the wire surface provides an opportunity to focus the electron beam, leading to an area of maximum density and correspondingly maximum energy density of hot electrons [19, 25]. We performed simulations for a target consisting of parallel wires. The target consists of carbon wires with diameter of \( s \geq 40 \text{ nm} \) and a separation distance \( \geq 100 \text{ nm} \). The laser with a wavelength of 0.8 \( \mu \text{m} \), a focal spot size of \( d_l = 4 \mu \text{m} \), a pulse duration (Gaussian pulse) of 45 fs and an intensity of \( 3 \times 10^{19} \text{ W cm}^{-2} \) illuminates the target such that the laser axis is parallel to the nanowires. We observe that in case of parallel wires, optimum conditions for number and energy of the hot electrons are fulfilled, if the target diameter is of the order of the skin layer \( l_s \) and the distance between the wires is more than the Debye radius of the hot electrons. The simulation shows that in such targets a bunch of hot electrons can be created which propagates deeply into the target. In figure 7(a) the propagation of the relativistic electron beam along the target wires and in figure 7(b) the propagation of the transverse component of the electric field is shown. Apparently, the laser field does not propagate along the wires but it transforms into a surface wave propagating along the wires with a velocity close to the speed of light \( c \) together with the relativistic electrons, supporting its propagation.

These results are in contrast to dense targets, where only very few hot electrons can reach comparable distances. The physical reason of extended hot electron propagation along the wires is the generation of a return current of cold electrons in the wires: hot electrons which partly leave the wire cause a positive charge which attracts cold electrons in the surrounding wire volume. The density of the return current in the wire is larger than the hot electron current density in the wire, but both currents compensate each other if the hot electrons propagating in the vacuum inter-space are added. Thus, the hot electron bunch can propagate along the wire over a remarkable distance. The analysis of the angular distribution of the hot electrons shows that the maximum electron density occurs within an angle of \( 10^\circ \). This means, while the electrons are propagating along the wire at the speed close to \( c \) they undergo transverse oscillations (3D spiral-like motion) in the electrostatic field of the wire which is termed ‘Debye wrapping’ of the propagating electron bunch.

**Experimental observations**

Experiments on intense laser pulse interaction with similar targets have been performed at the MBI High Field Ti:Sapph laser. It delivers \( 1.3 \) J on target with pulse durations of \( 30–35 \) fs. Contrast enhancement by an XPW [20] frontend leads to a peak to ASE contrast of \( \sim 10^{11} \). An ultra-short (<100 fs) pre-pulse at a level of \( 10^{-7} \) peak intensity precedes the main pulse by approximately 8 ns and is due to residual surface reflection.
inside the laser chain. Other experiments [23] showed that this pre-pulse does not destroy or significantly heat a 30 nm thin polymer foil. The laser is focused by F/2.5 off-axis parabola mirror to a focal FHWM size of 4 μm, yielding a peak intensity of $6 \times 10^{19}$ W/cm². We focused the laser on free-standing 5 μm thick Ti foil with different nano-structures. In the first case (see figure 7(a)), a brush target of irregular structure was engineered by chemical bath deposition similar to [21]. In the second case, we used ripple targets, which have been prepared by laser induced periodic surface structuring (LIPSS) [14, 22] using a separate kHz fs-laser.

The experimentally obtained energy spectral distribution of protons and carbon ions are displayed in figures 8(a) and 9(a) for interaction with plane foil and nanostructured foil targets. It is clearly visible that the ion number is significantly enhanced for both structures in comparison to the plain foil case. The brush structure furthermore shows a small enhancement in the energy cutoff of C⁴⁺, while the cut-off energy remains almost unchanged for the ripple structure.

In our simulations we used Ti^{15+} foils where the rear sides were covered by the contamination layer of 40 nm thickness consisted from the mixture of three ion sorts H⁺, C⁺⁺, Ti⁺⁺⁺ at different proportions. In the case of the proportion: 0.1 : 0.2 : 1 we obtain the calculated distributions close to the experimental ones. Thus the distribution of absorbed energy in a relief target of different ion sorts depends on its densities.

In the experiment we observe an increase of integrated carbon ion energy from the plane to the ripple target by a factor of ~1.4. In the simulation we calculated the integrated energy of C⁺⁺ ions correspondingly with $E_{\text{tot}} = \int E(dN/dE) \text{d}E$ and found that the ratio is $E_{\text{tot}}(\text{ripple})/E_{\text{tot}}(\text{plain}) \approx 2.3$. The small mismatch is due to the simulation geometry.

**Conclusion**

In summary, we have shown that a nano-relief structure with optimum parameters can significantly increase the short-pulse laser absorption, which is related to an enhancement of electron movement between relief ledges. A developed analytical model permits the approximation of the absorption enhancement in dependence on the relief parameters. The optimum relief height decreases with increasing laser intensity. But for the higher intensities the efficiency of such target stays high because here a transition to another dominant acceleration mechanism—the ‘laser piston’ [15] sets in. In the considered cases, degradation of the structure by a laser prepulse is crucial and a very high temporal contrast of the laser-pulse is required. The demand on an ultrashort laser pulse duration is relaxed, since after destruction of the original relief a secondary dynamic structure of ion density appears. Thus, strong absorption related to a plasma relief persists for a relatively long time of several hundred fs. This opens further options for the applications of structured targets. As an alternative approach a nano-wire target (very high relief) has been investigated permitting almost total absorption of the high intensity laser pulse, which subsequently is converted into a bunch of relativistic electron beams. They can propagate along the nano-wires for a long distance (tens of mm). As these electron bunches can follow a denting of the wire they can be focused into a small spot. The presented experimental results in particular confirm the theoretically predicted enhancement of ion numbers from laser irradiated nano-structured targets. Further investigation concerning achievable maximum ion energies by additional target parameters (thickness, structure size) remains necessary.

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